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ACTIVE EXPERIMENTS IN SPACE: PAST, PRESENT, AND FUTURE

EDITED BY: Gian Luca Delzanno, Joseph Eric Borovsky and Evgeny V. Mishin

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ACTIVE EXPERIMENTS IN SPACE: PAST, PRESENT, AND FUTURE

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Editorial: Active Experiments in Space: Past, Present, and Future

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Keywords: active space experiments, plasma physics, magnetospheres, ionosphere, laboratory astrophysics, space physics

Editorial on the Research Topic

Active Experiments in Space: Past, Present, and Future

Between 1958 and 1962 the United States and the Soviet Union performed several nuclear detonation tests in the atmosphere, including the Starfish Prime event which involved a 1.4 Mt explosion at 400 km altitude over Johnston island on July 9, 1962 (Gombosi et al., 2017). These tests can be considered as the beginning of active experiments in space (i.e., experiments that deliberately perturb the local environment). They demonstrated the potential destructive power of high-altitude nuclear explosions, both in terms of the resulting electromagnetic pulse as well as for the creation of a potentially long-lasting artificial radiation belt from the radioactive fission debris. For instance, one of the unintended consequences of Starfish Prime was to cripple at least seven spacecraft in low-Earth orbit (LEO), about a third of the LEO spacecraft of the time (Gombosi et al., 2017).

At about the same time, the fundamental discovery of the Earth's radiation belts by Van Allen and his team (Van Allen and Frank, 1959 and references therein) indicated how harsh the space environment could be for spacecraft and astronauts as well as how little we knew about it. Following the impetus of the Space Age, active space experiments flourished with the goals of (1) probing basic plasma physics phenomena, (2) elucidating aspects of magnetospheric and ionospheric physics, and (3) understanding how to control the effects of the environment on space assets. Bombs, beams, heaters, releases, chemical dumps, plasma plumes, tethers, antennas, voltages are examples of active experiments spanning several decades of research.

Six decades later the US active space experiment program has changed dramatically. The number of space-based experiments has seen a steep decline, supplanted by ground-based experiments that study the heating and modification of the ionosphere induced by powerful transmitters, such as the facilities of the High-Frequency Active Auroral Research Program (HAARP) and at Arecibo. This decline can be attributed to several reasons, summarized by the fact that the “low-hanging fruits” had already been collected and much more is known today about the space environment, that space flight became more bureaucratic and more risk-adverse, and budgetary pressures (Delzanno and Borovsky, 2018).

Yet, there are many reasons to be optimistic about the future of active experiments in space. There are new scientific and national-security drivers that demand new active space experiments. One example involves magnetosphere-ionosphere coupling, where a high-power electron beam could be used for magnetic field line mapping and connect phenomena occurring in the distant magnetosphere with their image in the ionosphere (National Research Council, 2012). Another example concerns radiation belt remediation, where the fluxes of an artificial radiation belt created by a high-altitude nuclear explosion could be substantially reduced by space-based injection of electromagnetic plasma waves with the objective of protecting critical space assets. Furthermore, there are new maturing technologies (metamaterials, compact relativistic

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accelerators, antennas constructed of superparamagnetic nanoparticles, cube-satellites, ...), ever better diagnostics and data-gathering capabilities, and new computational tools that can support the design and interpretation of active experiments like never before. Indeed, this optimism was conveyed by the 65 participants of the Workshop “Active Experiments in Space: Past, Present and Future” who gathered in Santa Fe, New Mexico, in September 2017 (see the link <http://www.cvent.com/events/active-experiments-in-space-past-present-and-future/event-summary-73675ac6ba5745d48d181933c4783454.aspx?dvce=1> for a list of talks presented at the workshop) and is echoed in this Frontiers special issue with the same name.

This special issue was designed to connect the past, the present, and the future of active experiments in space and serve as a reference for the community. It involves several review articles of past active experiments discussing chemical releases (Haerendel) and diamagnetic cavities (Winske et al.), artificial aurora experiments (Mishin), the APEX electron-beam experiments (Prech et al.), and an overview of active experiments involving the Los Alamos National Laboratory (Pongratz). A review of more recent active experiments focusing on electromagnetic wave injection and wave-particle interaction physics is given by Golkowski et al.. A review on the potential use of electron beams to solve outstanding problems in space physics is presented by Sanchez et al., while a review on the future of active experiments is presented by Borovsky and Delzanno. The special issue also contains several “Original Research” articles. A major focus is on the research associated with the use of high-power electron beams for space applications. These articles include a discussion of the development of new, compact, relativistic electron accelerators (Lewellen et al.), a tether-based spacecraft charging mitigation scheme (Marchand and Delzanno), the evolution of a relativistic electron beam for magnetic-field line mapping in near-Earth space (Powis et al.) and how magnetic-field-line curvature affects the ionospheric accessibility of the electron beam (Willard et al.), a method for measuring the local magnetic-field-line curvature in the inner magnetosphere with a variable-energy electron beam (Willard et al.), and the atmospheric signatures created by relativistic electron beams (Marshall et al.). New ionospheric experiments involving very-long-distance propagation of high-frequency waves are discussed by

Yampolski et al. A simulation study of the effect of plasma releases on the equatorial spread F is presented by Zawdie et al., while the use of plasma releases to enhance energetic neutral atom imaging is discussed by Scime and Keesee. Active experiments for planetary missions are discussed by Gilet et al. for active probes and by Voshchepynets et al. for sounding radar operations.

In the words of Nobel laureate Hannes Alfvén (Alfvén, 1970): “The center of gravity of the physical sciences is always moving. Every new discovery displaces the interest and the emphasis. Equally important is that new technological developments open new fields for scientific investigation. To a considerable extent the way science takes depends on the construction of new instruments as is evident from the history of science.” This is certainly true for active experiments and an exciting new season of active experiments in space awaits. *Ad maiora!*

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All authors listed have made a substantial, direct and intellectual contribution to the work, and approved it for publication.

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Tethered Capacitor Charge Mitigation in Electron Beam Experiments

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Energetic electron beams have been proposed for tracing magnetic field lines from the magnetosphere down to the ionosphere, in active experiments aimed at diagnosing mechanisms at play in the coupling between magnetosphere and ionosphere. It is recognized however that in the absence of an efficient mitigation technique, this approach would lead to unacceptably large spacecraft charging and positive potential buildup, which would result in environmental hazard for the spacecraft. This problem would be particularly acute in low density regions of the magnetosphere of interest in the study of magnetic field reconnection and substorm dynamics. A solution to this predicament could consist of creating a plasma contactor whereby a gas puff would be ionized, leading to the evacuation of positive charges and collection of cold electrons, thus compensating for the charges lost in the electron beam. A possible alternative is presented here, which consists of attaching a large passive conducting surface to the spacecraft, a “tethered capacitor”, from which negative charges would be drawn to compensate for those lost from the beam. This capacitor would then charge to a large positive potential, leaving the spacecraft and electron gun at a lower, acceptable positive potential. The tethered capacitor could have a relatively small mass; consisting only of a thin conducting surface that would be “inflated” as a result of repulsive electrostatic forces. This charge mitigation concept, as applied to active electron beam experiments, is explored using three dimensional particle-in-cell (PIC) simulations from which scaling laws can be inferred for the spacecraft and tethered capacitor potentials under proposed electron beam operations.

Keywords: charge mitigation, charge collection enhancement, tenuous magnetospheric plasma, magnetosphere-ionosphere coupling, spacecraft charging, electron beam, tethered capacitor

1. INTRODUCTION

Electron guns have been used on several satellites to perform a variety of active experiments, including controlling the floating potential (Whipple and Olsen, 1980; Koons and Cohen, 1982; Pedersen et al., 1984), and diagnosing distant parts of the magnetosphere along magnetic field lines (Hendrickson et al., 1975; Wilhelm et al., 1980; Winckler, 1980, 1992; Nemzek et al., 1992; Grandal and North, 2012). An interesting proposal to elucidate the long-standing problem of magnetosphere-ionosphere coupling and, more generally, of what creates auroral forms, involves the injection of energetic electrons along field lines (i.e., in the loss cone), from satellites in near geostationary orbits. By monitoring light emitted by atoms excited by these electrons entering

the upper atmosphere, this approach would map unambiguously the ionospheric footpoint of the magnetospheric spacecraft. Combined with suitable measurements of the magnetospheric and ionospheric state, it would determine what magnetospheric conditions are related to the wide variety of auroral forms. The approach considered here consists of emitting 1 MeV electron pulses along magnetic field lines from a satellite in elliptic near geostationary orbit, sampling the magnetosphere at altitudes between 30,000 and 50,000 km. The orbit inclination and argument of the perigee would be such as to maximize the time spent over the TReX array in Canada (Donovan, 2015a,b). Electrons would be injected in a succession of 0.5 s pulses, each comprising 100, 10 mA minipulses of duration 0.5 ms followed by a 4.5 ms rest period. Following each 0.5 s pulse, there would be a half second of rest to allow for batteries to recharge, and the process would be repeated. The average electron current emitted over the 0.5 s period would then be 1 mA. This pulsed operation of the electron beam would enable the detection of the beam spot in the atmosphere through blinking against a bright auroral background. One technical difficulty is that owing to the low plasma density in the region of interest ($n_e \sim 10^6 \text{ m}^{-3}$), neutralization from collecting background electrons alone would not be sufficient to prevent the spacecraft from reaching unacceptably large positive potentials. Indeed, with space parameters of interest here, assuming balance between background electron collection and the 1 mA average current emitted during the first 0.5 s of each pulse, a simple orbital-motion-limited (OML) (Mott-Smith and Langmuir, 1926; Sudit and Woods, 1994; Allen et al., 2000) estimate of the floating potential leads to a value of the order 45 kV. Not only these values are unacceptably large for spacecraft safety but they would also cause energetic electrons to strike spacecraft components, requiring careful consideration from the perspective of internal charging. A promising solution proposed to avoid the buildup of such large potentials consists of using a “plasma contactor” whereby a relatively dense plasma cloud would be created near the spacecraft. As the spacecraft potential increases, electrons from the plasma would be attracted to the satellite while ions would be repelled (Delzanno et al., 2015; Lucco Castello et al., 2018), thus maintaining the floating potential at a lower value through emission of significant ion currents by the contactor cloud. Plasma contactors or hollow cathodes have proven their effectiveness in laboratory (Stenzel and Urrutia, 1990; Williams and Wilbur, 1990) and space (Olsen, 1985; Patterson et al., 1993; Katz et al., 1994; Comfort et al., 1998; Safránková et al., 2002) for controlling satellite floating potentials. In addition, ion beam emission (Schmidt et al., 1995; Torkar et al., 2001) and the release of neutral gas (Gilchrist et al., 1990) have also been used to mitigate satellite charging in space. The alternative considered here consists of attaching a large tethered passive surface capacitor, from which electrons would be drawn as needed by the active spacecraft in order to maintain its floating potential to an acceptable level. This capacitor could consist of an electrostatically “inflatable” conducting foil held by a boom at a fixed distance from the satellite. During the emission of electron pulses, as negative current would be drawn from the capacitor, its potential would increase to large positive values until on average,

the capacitor and the spacecraft collect sufficient background electrons to balance lost electron beam charges. The feasibility of this concept is assessed quantitatively in the following sections.

In section 2, we first present a simple estimate of the conditions required for current balance for the spacecraft-capacitor system when collected currents and voltages are assumed to be constant. These results are then used to derive an empirical scaling law for the current collected by the capacitor under various operational scenarios. The following section considers the time dependence of the voltages and collected currents when the electron beam is in operation. This is then followed by section 4 in which a simple analytic model is presented to estimate the peak potential and collected currents at the end of 0.5 ms of beam emission. A summary and conclusions are presented in section 5

2. STEADY STATE CURRENT BALANCE

Our goal is to find a satellite-tethered-capacitor configuration that satisfies current balance during the 0.5 s when the electron accelerator is operated, such that the satellite and capacitor potentials remain at an acceptable level. To start with, we consider steady state conditions, in which the spacecraft and capacitor are at fixed potentials, and calculate the currents collected under different conditions. For the purpose of the estimate, a simple geometry is assumed for the assembly. It consists of a cubic satellite bus with side length $l_{SC} = 2 \text{ m}$, a thin tether of length l_{teth} to which a spherical conducting sphere of radius a_{cap} is attached, as illustrated in **Figure 1**. The spacecraft surface is assumed to be equipotential at voltage $V_{SC} = 1 \text{ kV}$, while the tether and capacitor are assumed to be at the same voltage V_{cap} . For simplicity, a pure fully ionized non-drifting electron-proton plasma background is assumed, and the magnetic field is assumed to be zero. Our calculations are based on fully kinetic Particle In Cell (PIC) simulations made with PTetra (Marchand, 2012; Marchand and Resendiz Lira, 2017), an explicit electrostatic code which uses unstructured tetrahedral meshes to discretize the simulation domain. In the simulations we consider tether lengths l_{teth} varying from 5 to 10 m, and capacitor voltages V_{cap} from 10 to 25 kV. In all cases, including the time dependent solution considered in section 3, the emitted electron beam is not included in the simulations. The charge lost in the beam is accounted for indirectly by subtracting the emitted charges from the capacitor at each simulation time step. That is, the electron gun mounted on the spacecraft is assumed to draw all its current from the capacitor. The simulation domain is discretized with a mesh consisting of approximately 3.5 million tetrahedra and 600,000 vertices. Approximately 200 million macro-particles are used to represent electrons and ions in the PIC simulations. In all cases the simulation domain is delimited by a spherical boundary of radius $R_{boundary} = 100 \text{ m}$ where the potential is assumed to be zero. The radius of the outer boundary must be sufficiently large for the incoming electron flux at the boundary to be significantly larger than the average 1 mA that must be collected by the spacecraft-capacitor assembly. If the outer boundary is not far enough and the incoming flux is <1

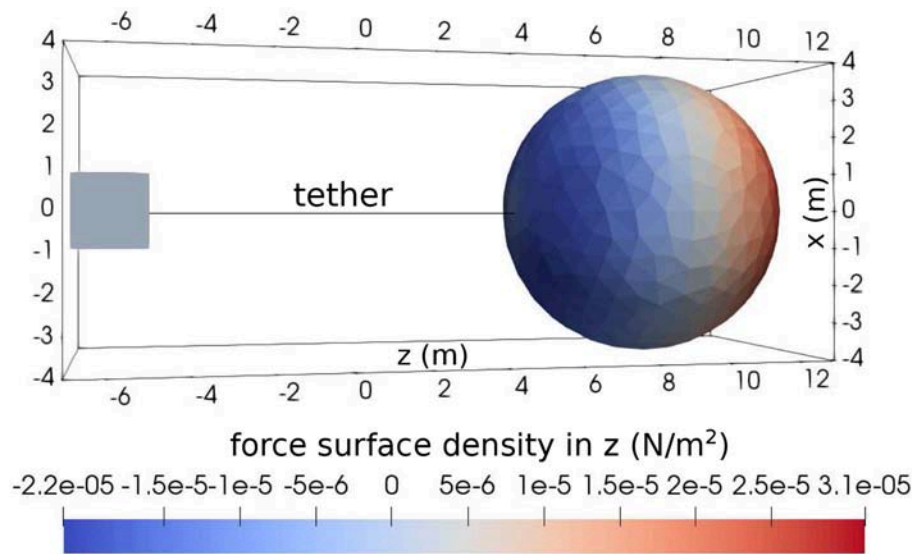


FIGURE 1 | Illustration of the spacecraft, tether, and spherical capacitor geometry assumed in the simulations. The colors on the sphere show the component of the electric force surface density along z , computed at $t = 5.5$ ms in the time dependent simulation.

TABLE 1 | Physical and simulation parameters considered in the simulations.

ion species	100% H^+	l_{SC}	2 m
$n_e = n_{H^+} = n$	$10^6 m^{-3}$	l_{teth}	5, 7, 10 m
$T_e = T_i = T$	1 keV	a_{cap}	2, 3, 4 m
λ_D	166 m	$R_{boundary}$	100 m

mA, the entire simulation becomes nearly depleted of electrons, the spacecraft potential increases without bounds, and no steady state solution can be found. Assuming a background Maxwellian distribution, in order for the electron flux at the outer boundary to match the required current, the following condition must be satisfied (Laframboise, 1966):

$$I_0 = 10^{-3} A = n_e e R_{boundary}^2 \sqrt{\frac{8\pi kT}{m_e}}, \quad (1)$$

where k is the Boltzmann constant, and m_e is the electron mass. With the space-plasma parameters considered here this gives a radius of approximately 10 m. Different radii have been considered including 50, 75, and 100 m. Only the latter has been used throughout as being sufficiently far from the spacecraft-capacitor assembly. A summary of the plasma and simulation parameters is given in **Table 1**.

Collected currents were computed with different values of tether lengths l_{teth} , capacitor radius a_{cap} and capacitor potentials V_{cap} . **Table 2** lists the currents collected by the spacecraft bus, the capacitor, and the total collected current computed with the different configurations considered. The range of parameters was selected so as to yield a total collected current close to the $I_{total} = 1$ mA objective, and from which a simple analytic fit can be

TABLE 2 | Steady state currents collected by the spacecraft and the tethered capacitor for the different configurations considered.

l_{teth} (m)	a_{cap} (m)	V_{cap} (kV)	I_{SC} (mA)	I_{cap} (mA)	I_{total} (mA)
5	2	15	-0.039	-0.645	-0.684
5	2	20	-0.041	-0.852	-0.893
5	2	25	-0.042	-1.066	-1.108
7	3	8	-0.036	-0.799	-0.835
7	3	10	-0.038	-0.984	-1.022
10	2	15	-0.038	-0.713	-0.751
10	2	20	-0.039	-0.945	-0.984
10	2	25	-0.038	-1.180	-1.218
10	4	15	-0.038	-2.594	-2.632
10	4	20	-0.038	-3.423	-3.461
10	4	25	-0.038	-4.240	-4.278

In all cases the spacecraft potential is $V_{SC} = 1$ kV.

constructed. The table shows that in all cases considered, most of the collected current is from the spherical capacitor. Thus guided by the OML expression for the electron current collected by a positive conducting sphere, we fitted the total collected current with

$$I_{total} = -en_e a_{cap}^2 \sqrt{\frac{8\pi kT}{m_e}} \times (\alpha + \beta l_{teth}) \left(1 + \frac{V}{kT_{eV}}\right)^\gamma, \quad (2)$$

where e is the unit charge and T_{eV} is the temperature in eV. In Equation 2 fitting parameters α , β , and γ are set to best approximate simulation results. For an isolated sphere, idealized OML theory would predict 1, 0, and 1 for α , β , and γ , respectively. Deviations from these values are allowed to account for the

presence of the nearby spacecraft, the length, and bias voltage of the tether. This expression is fitted to the data in **Table 2** so as to minimize the maximum relative error. The result, $\alpha = 0.983$, $\beta = 0.050$, and $\gamma = 1.000$, fits all data points with a maximum relative error of 6.6%. This analytic fit can also be used to find conditions under which a given current will be collected at steady state, for satellite and tethered capacitor parameters in the range of those considered in **Table 2**. For example, a 7 m tether with a 3 m radius capacitor biased to 10 kV should be adequate to balance the average 1 mA current lost while emitting electron pulses on the “Magnetosphere-Ionosphere Observatory (MIO)” (Borovsky, 2002; Delzanno et al., 2016). From the fit and **Table 2**, it is also clear that different electron currents could be collected from background plasma, with modified configuration parameters.

3. TIME-DEPENDENT SOLUTION

The results presented so far have focused on currents collected assuming constant spacecraft and capacitor voltages. From **Table 2** and the analytic fit 2 it appears that under these conditions it should be possible for the spacecraft-capacitor assembly to collect the 1 mA current required to balance the average current lost during the 5 ms period of a minipulse (i.e., 0.5 ms of electron gun operation at 10 mA followed by 4.5 ms without emission). The steady state currents for given voltages would be applicable to the actual experiment provided that the spacecraft and capacitor voltages remain approximately constant during a full 5.0 ms of a minipulse. This in turn would be valid if the relaxation time of the system were larger than 5.0 ms; that is, if the voltages didn’t change significantly during the 4.5 ms beam rest period. However, simulations indicate otherwise. The relaxation times for the voltage and collected currents are found to be of order 2 ms, which implies that while the estimates obtained assuming steady state potentials are indicative of how collected currents relate to voltages under the assumed geometries, they cannot be seen as quantitatively accurate. This is apparent in **Figures 2, 3** for the time evolution of the voltages and collected currents, assuming a 10 m tether and a 4 m spherical capacitor. These results were obtained assuming a uniform plasma and zero potential on all components at time $t = 0$. In the simulation, the full current of the electron beam is assumed to come from the sphere. Any charging of the spacecraft is from the impact of surrounding plasma particles. During the first 0.5 ms when electrons are being fired, the capacitor and spacecraft voltages increase almost linearly to 12.3 and 4.66 kV, respectively. This is accompanied by a nearly linear drop in the collected current reaching -2.23 and -0.12 mA for the sphere and spacecraft, respectively. In the following 4.5 ms rest period both voltages and currents relax quasi-exponentially toward an equilibrium. The first 0.5 ms of the following minipulse was also simulated to ascertain whether differences in the initial condition would result in significant changes in the time evolution of the system. At time 5.5 ms, the spacecraft voltage is 3.47 kV and that of the capacitor, 13.0 kV, while the collected currents are -0.10 and -2.3 mA, respectively. These values being relatively close to the ones found at $t = 0.5$ ms, we conclude that the time evolution

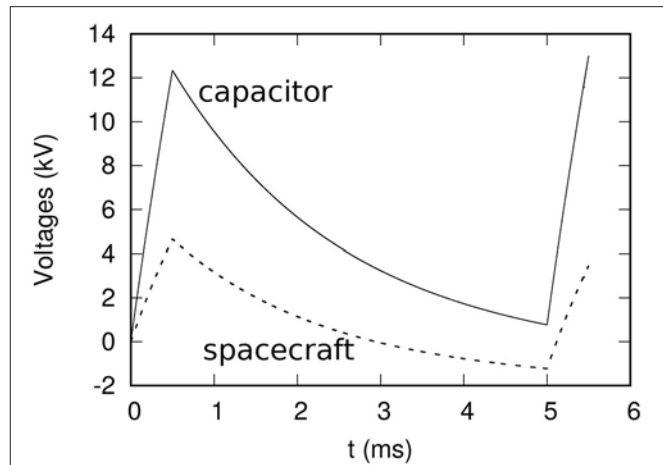


FIGURE 2 | Spacecraft and capacitor voltages over a 5.5 ms time period.

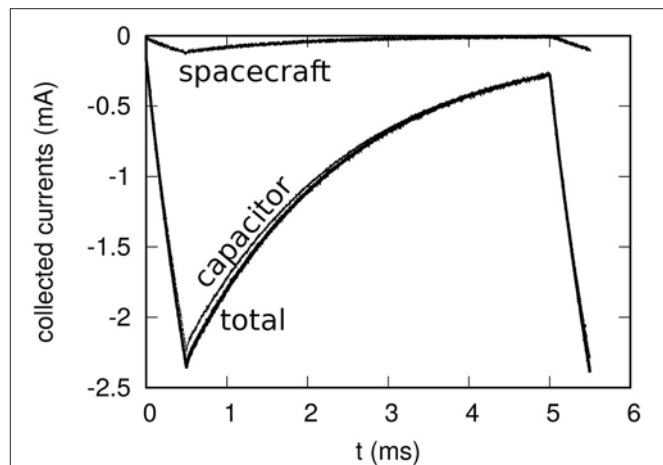


FIGURE 3 | Collected current by the spacecraft, the capacitor and the sum of the two, over a 5.5 ms time period.

of the potentials and collected currents calculated in the first 5.0 ms provide a good representation of what would be found within subsequent minipulses. It follows from the rapid decay, occurring on a time scale of order 2 ms, that both satellite and capacitor would equilibrate with space environment within 20–50 ms. This is longer than the 4.5 ms separating successive minipulses, but sufficiently short for an equilibrium to be reached during the 0.5 s rest period separating different pulses.

It is noteworthy that toward the end of the 5 ms cycle, the spacecraft potential becomes negative. This is due in part to the decrease in the capacitor potential, and to the collection of a fraction of the electrons attracted to the capacitor. Finally, the force between the spacecraft bus and capacitor has been calculated by integrating the force surface density; that is, the force per unit surface area, over both surfaces, as

$$\vec{F} = \iint_S ds \frac{1}{2} \epsilon_0 E^2 \hat{n}, \quad (3)$$

where the integration is carried out over the entire surface of a given object, ϵ_0 is the free space permittivity, E is the magnitude of the electric field at the surface, and $\hat{n}$ is the unit vector perpendicular to the surface pointing outward. The z component of the force surface density exerted on the spherical capacitor computed at time 5.5 ms is illustrated in **Figure 1**. The figure shows an approximate $\cos(\theta)$ profile, where θ is the angle between the direction along the tether and the radial position on the sphere, with an offset of approximately $4.5 \times 10^{-6} \text{ Nm}^{-2}$. The computed net force along the z direction is $F_z \simeq 3.4 \times 10^{-4} \text{ N}$, corresponding to a repulsive force with respect to the satellite. The force between the spacecraft and the sphere, while small, would likely become attractive during part of a 5.0 ms period; in particular when the spacecraft becomes negative while the sphere remains positive. If the average force were to be attractive, a possible solution would be to use a tether with sufficient rigidity to keep spacecraft and capacitor apart.

4. ANALYTIC APPROXIMATION

The time dependent results just considered suggest a simple analytic model to capture the essential of the system response to the emission of electron beam pulses. Assuming that the potential at the spacecraft is due mainly to the influence of the nearby capacitor, the spacecraft potential should be

$$V_{SC} \sim V_{cap} \frac{a_{cap}}{l_{teth} + a_{cap} + l_{SC}/2}. \quad (4)$$

Considering the large Debye length in this environment, the capacitor voltage is approximated as for a sphere in vacuum:

$$V_{cap} \simeq \frac{Q}{4\pi\epsilon_0 a_{cap}}, \quad (5)$$

where, assuming zero initial charge,

$$Q = \left(-I_{beam} + \frac{I_{plasma}(t)}{2} \right) t, \quad (6)$$

is the negative of the charge emitted in the beam in time t , plus the charge collected from incident background electrons. In this expression, use is made of the fact that, during the 0.5 ms period when the electron beam is fired, $I_{plasma}(t)$ increases approximately linearly with time. Omitting the explicit time dependence in I_{plasma} for brevity and assuming $V > 0$, we then approximate I_{plasma} with the OML theory, which yields

$$\begin{aligned} I_{plasma} &\simeq -ena_{cap}^2 \sqrt{\frac{8\pi kT}{m_e}} \left(1 + \frac{eV_{cap}}{kT} \right) \\ &\simeq -ena_{cap}^2 \sqrt{\frac{8\pi kT}{m_e}} \left(1 - \frac{e(I_{beam} - I_{plasma}/2)t}{4\pi\epsilon_0 a_{cap} kT} \right) \\ &\simeq \frac{-ena_{cap}^2 \sqrt{\frac{8\pi kT}{m_e}} \left(1 - \frac{eI_{beam}t}{4\pi\epsilon_0 a_{cap} kT} \right)}{1 + \frac{e^2 na_{cap}}{2} \sqrt{\frac{8\pi kT}{m_e}} \frac{t}{4\pi\epsilon_0 kT}}. \end{aligned} \quad (7)$$

Substituting 7 in 6, and evaluating the numerical values of physical constants, equation 5 becomes

$$V \sim 9 \times 10^9 a_{cap}^{-1} \left[-I_{beam} + \frac{I_{beam}t - 10^{-7} a T_{keV}}{t + 2 \times 10^4 n^{-1} a^{-1} T_{keV}^{1/2}} \right] t, \quad (8)$$

where V , a_{cap} , I_{beam} , and t are in SI units, and T_{keV} is the temperature in keV. With the parameters considered in the time dependent simulation, $a = 4 \text{ m}$, $I_{beam} = -0.01 \text{ A}$, and $T = 1 \text{ keV}$, the predicted capacitor potential from 8 at time $t = 5 \times 10^{-4} \text{ s}$ is $V \sim 10 \text{ kV}$. Using this estimate for V in the top part of Eq. 7, we find $I_{plasma} \sim 1.9 \text{ mA}$. Finally, the spacecraft potential estimated with 8 at $t = 0.5 \text{ ms}$ is $V_{SC} \sim 2.67 \text{ kV}$. Compared with computed values, these estimates are of course approximate, with relative errors ranging from $\sim 20\%$ for the predicted capacitor voltage and collected current from background plasma, to $\sim 50\%$ for the spacecraft potential. While approximate, this simple model should be useful in exploring parameter space and finding optimal conditions for the proposed approach to be applied.

5. SUMMARY AND CONCLUSION

An alternative to plasma contactors has been presented as a possible means of mitigation for spacecraft charging occurring when electron beams are emitted in a tenuous plasma. The solution consists of attaching a capacitor with a large conducting surface area to a spacecraft from which current can be drawn to compensate for the current injected by the electron beam. As a result, the capacitor becomes strongly positively charged, thus reducing positive increases in the spacecraft potential, and mitigating possibly adverse effects. For the parameters considered here, with a collecting surface area larger than that of the spacecraft, the conducting capacitor can collect sufficient electron current from background plasma to balance an average electron beam current of order 1 mA, while maintaining a potential of order 10 kV. Assuming constant voltages for the spacecraft and capacitor, scaling laws were derived for the collected current which suggest that it might be practical to apply this approach to mitigate spacecraft charging in active electron beam emission experiments. A simulation has also been made to explore the time evolution of the voltages and collected currents. Three time periods have been considered, including the 0.5 ms of beam emission, followed by a 4.5 ms when the electron beam is turned off, and a subsequent 0.5 ms of beam emission. Voltages and collected currents exhibit significant variations in time, with extrema at the end of 0.5 ms period of beam emission, with an approximate exponential decay during the following 4.5 ms rest period. The average values found for the collected currents, however, are consistent with results obtained in the steady state approximation. A simple analytic model was constructed that can capture the dependence of the tethered capacitor voltage and collected current as a function of time and other physical parameters. Predictions made with this model can only be viewed as approximate, but they should prove useful, in conjunction with kinetic simulations, to explore and optimize system parameters.

In this paper, we have focused on the conditions where the spacecraft charges to values of order of 1 kV. This is acceptable

from the perspective of the emission of the electron beam (since the beam energy is much larger than 1 kV) and for spacecraft safety (provided that the spacecraft platform is designed to avoid differential charging and the risk of electrostatic discharges). However, kV potentials can perturb local plasma measurements, meaning that one would then have reliable measurements only in-between (0.5 s) pulses. One could increase the tether length and/or the capacitor size to decrease the spacecraft potential. Furthermore, in a real mission the main spacecraft carrying the electron gun could be accompanied by daughter spacecraft sufficiently far not to be affected by perturbations resulting from the beam emission, and from which plasma environment parameters and their gradients would be measured (Borovsky, 2002; Dors et al., this issue).

Clearly, for this approach to be practical, it would be necessary for on-board electronics to draw negative current from the capacitor when it is at voltages well above that of the spacecraft. Assuming a 10 kV potential difference, the energy needed to extract a charge of $-10 \text{ mA} \times 0.5 \text{ ms}$ from the spheres should be of order 0.05 J, which is much less than the $10 \text{ mA} \times 0.5 \text{ ms} \times 1 \text{ MV} = 5 \text{ J}$ energy in each electron pulse. We recall that several simplifying assumptions have been made in this work. The spacecraft, tether, and capacitor have been somewhat idealized, and several physical effects such as photo-electron emission and secondary electron emission have been neglected. Considering the typically large positive voltages of the assembly and the low energy of photo or secondary electrons (a few eV) however, any emitted electron would almost certainly be accelerated back to the emitting surface. One point to consider is the possibility that a fraction of the background incident energetic electrons, with thermal energy 1 keV plus the energy of $\sim 10 \text{ kV}$ gained when reaching the capacitor, would go through it if the material that the capacitor is made of does not have sufficient stopping power.

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In this study, a simple sphere was considered, with the suggestion that it could consist of an electrostatically “inflatable” foil. In order to ensure that incident energetic electrons be captured by the capacitor, however, other geometries with adequate material properties should be considered in future studies.

DATA AVAILABILITY STATEMENT

The datasets analyzed in this study are contained in the manuscript itself.

AUTHOR CONTRIBUTIONS

RM was mainly responsible for carrying out the simulations. RM and GD contributed though discussions and in the interpretation of the results.

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History of Los Alamos Participation in Active Experiments in Space

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Los Alamos has a long history of participation in active experiments in space beginning with the Teak nuclear test in 1958. Above-ground nuclear testing stopped in 1962 because of the Partial Test Ban Treaty, and a program of non-nuclear chemical release experiments began in 1968. Los Alamos has participated in nearly 100 non-nuclear experiments in space, the last being the NASA-sponsored strontium and europium doped barium thermite releases in the Arecibo beam in July of 1992. The rationale for these experiments ranged from studying basic plasma processes such as gradient-driven structuring and velocity-space instabilities to illuminating the convection of plasmas in the ionosphere and polar cap to ionospheric depletion experiments to the B.E.A.R. 1-MeV neutral particle beam (NPB) test in 1989. This report reviews the objectives, techniques and diagnostics of Los Alamos participation in active experiments in space.

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INTRODUCTION

A request to give a talk at the “Active Experiments in Space: Past, Present and Future” conference in Santa Fe in September 2017 motivated the research behind this report. We were asked to summarize LANL involvement in active experiments (<http://www.cvent.com/events/active-experiments-in-space-past-present-and-future/event-summary-73675ac6ba5745d48d181933c4783454.aspx>). The previous summary paper on LANL-related active experiments is out of date [1]. At Los Alamos active experiments have been somewhat of a stepchild to all the wonderful discoveries made by LANL satellites beginning with the VELA satellites. However, active experiments play a complementary role to the exploration satellites. The hypotheses testing active experiments play a “Galileo” role in space science while the instrumented satellites and rockets play a “Christopher Columbus” discovery role.

This report is a “history” and not a description of a scientific investigation. We begin with the people involved over the years because as you get older those are the associations you cherish. Then of course we need to identify the funding sources for our active experiments. Next, we present a “catalog” describing the timeline and locations of the many active experiments with Los Alamos involvement. Following this, we describe the various active experiment techniques and diagnostics we’ve employed. Next, we briefly describe the many objectives of these experiments. Lastly, we close with a few examples of our active experiments. An extensive bibliography provides additional detail and experiment results.

THE PEOPLE

For most of these experiments Los Alamos has not acted alone. Early on we partnered with our fellow AEC laboratory Sandia; they had the rockets and we had the shaped-charges and

cameras. Many suggestions for experiments and facilities and diagnostics came from our partners at the University of Alaska's Geophysical Institute. A number of the experiments studied auroral phenomena and our Canadian partners provided launch support and diagnostics. Over the years we have also partnered with EG&G, the Naval Research Laboratory, the Aerospace Corporation, the Lockheed Palo Alto Research Laboratories, Goddard Space Flight Center, the Max Planck Institute and others.

We've had the privilege of working with many great folks over the years. We begin with an "In Memoriam" tribute to Gene Wescott. Eugene Michael "Gene" Wescott (February 15, 1932–February 23, 2014) was an American scientist, artist, and traditional dancer. Wescott worked at the Geophysical Institute of the University of Alaska Fairbanks from 1958 to 2009. He was appointed Professor Emeritus of Geophysics, and had an extensive background of research in solid earth

geophysics and space physics. He was directly involved in auroral and magnetospheric electric field studies and plasma physics experiments using barium and calcium plasma rocket injections at Poker Flat Research Range. In the marriage between the Geophysical Institute and the AEC, Gene had the range and the ideas and the AEC had the rockets, the aircraft, and the shaped-charges.

The Los Alamos Scientific Laboratory, Los Alamos National Scientific Laboratory and Los Alamos National Laboratory (LASL/LANSL/LANL) involvement in many of the early active experiments originated in Group J-10. J-10 group leaders Herman Hoerlin, Milt Peek and Bob Jeffries supported the active experiments. Their successors in leading other LANL organizations including Doyle Evans, Don Cobb, and David Simons continued LANL support for active experiments in space. Their support was especially important in securing funding for active experiments. Many LANL staff members and technicians (Casey Stevens, Lois Dauelsberg, Hal Fishbine, Hal Dehaven, John Wolcott, Bob Carlos, Paul Bernhardt, and Gordon Smith to name just a few) were vital to our success. Special recognition goes to Mel Duran and the others manning an optical observatory in the dead of winter at Resolute Bay, NWT, Canada, during the Tordo and Periquito experiments.

It wasn't all cold weather either. Special thanks to the crew of the Beam Experiments Aboard a Rocket (B.E.A.R.) experiment conducted in summer heat of at the White Sands Missile Range in 1989. Of course, we conducted many experiments with our Sandia National Laboratory in the Hawaiian Islands. We also benefitted from the many interactions with our fellow Principal Investigators (PIs) in the NASA-sponsored Combined Release Radiation Effects Satellite (CRRES). A special acknowledgment to NASA's CRRES leader, David Reasoner.

TABLE 1 | Dates and locations of LANL-involved nuclear tests in space [2].

Event	Date	Latitude	Longitude	Altitude (km)
Teak	August 1, 1958	16.7	−167	76.8
Orange	August 12, 1958	16.7	−167	43
Argus I	August 27, 1958	−38.0	−10.9	
Argus II	August 30, 1958	−49.4	−8.7	
Argus III	September 6, 1958	−49.5	−10.45	
Starfish	July 9, 1962	16.7	−167	400
Checkmate	October 20, 1962	16.7	−167	
Bluegill	October 26, 1962	16.7	−167	
Kingfish	November 1, 1962	16.7	−167	
Tightrope	November 4, 1962	16.7	−167	

TABLE 2 | Dates and locations of LANL-involved thermite releases from orbit.

Event	Date	Time (UT)	Latitude	Longitude	Altitude (KM)
CAMEO 1	October 29, 1978	11:09:15	79.10	−112.43	968.4
CAMEO 2	October 29, 1978	11:09:55	77.74	−121.58	966.7
CAMEO 3	October 29, 1978	11:10:35	76.09	−129.33	965.3
CAMEO 4	October 29, 1978	11:11:15	74.31	−135.06	966.3
PEGSAT	April 5, 1990	19:10:00			
CRRES G-2	January 13, 1991	2:17:03	16.90	−103.1	6,180
CRRES G-7	January 13, 1991	7:05:00	8.00	−86.7	33,403
CRRES G-3	January 15, 1991	4:11:00	17.90	−97.5	15,063
CRRES G-4	January 16, 1991	6:25:00	−0.70	−53.8	23,977
CRRES G-5	January 18, 1991	5:20:00	6.60	−62.8	33,337
CRRES G-10	January 20, 1991	5:30:00	8.90	−75.6	33,179
CRRES G-6	February 12, 1991	4:15:00	4.90	−76.1	32,249
CRRES G-8	February 17, 1991	3:30:00	0.40	−58.1	33,553
CRRES G-1	July 13, 1991	35:25.5	17.80	−62.9	495
CRRES G-9	July 19, 1991	8:37:07	17.40	−62.8	441
CRRES G-11a	July 22, 1991	8:38:24	16.80	−60.3	411
CRRES G-11b	July 25, 1991	8:37:11	17.30	−69.5	478
CRRES G-12	August 12, 1991	9:31:20	9.10	−63.5	507

TABLE 3 | Dates and locations of LANL-involved thermite releases from rockets.

Event	Date	Time (UT)	Latitude	Longitude	Altitude (km)
APPLE	June 2, 1968	9:17:42	18.70	−66.8	196
DOGWOOD	June 12, 1968	9:10:42	18.70	−66.9	188
JUNIPER	March 4, 1969	32:23.0	64.53	−148.05	170
ELM	March 5, 1969	4:32:47	64.53	−148.05	170
IRONWOOD	March 11, 1969	14:10:47	64.53	−148.05	140
FIR	March 15, 1969	5:08:15	64.53	−148.05	165
GUM	March 19, 1969	5:19:15	64.53	−148.05	168
HEMLOCK	March 20, 1969	5:39:10	64.53	−148.05	176
ROADRUNNER	May 16, 1970	5:47:47	22.75	−159.752	234.39
SAPSUCKER	May 26, 1970	5:52:25	22.73	−159.738	211.613
TITMOUSE	June 6, 1970	5:55:25	21.41	−160.233	207.324
NUTMEG	January 16, 1971	23:34:40	30.63	−86.367	144
PLUM	January 20, 1971	23:47:05	30.74	−86.3	182
REDWOOD	January 26, 1971	23:52:09	30.88	−86.55	252
OLIVE	January 29, 1971	23:53:57	30.34	−86.13	352
SPRUCE	February 1, 1971	23:50:04	30.65	−86.55	184
CANUTO A	October 24, 1971	4:31:10	22.34	−160.485	201.93
CANUTO B	October 24, 1971	4:35:42	22.98	−162.149	201.89
DARDABASI A	November 8, 1971	4:28:10	22.42	−160.534	198.21
DARDABASI B	November 8, 1971	4:32:35	23.23	−162.271	197.9
ANNE	December 1, 1976	23:11:43			181
BETTY	February 26, 1977	52:27.2	29.70	−86.752	178.8
CAROLYN	March 2, 1977	54:10.5	29.69	−87.008	191.1
DIANNE	March 7, 1977	01:08.0	29.62	−86.662	85.5
ESTER	March 13, 1977	01:08.8	29.70	−86.807	189.2
FERN	March 14, 1977	46:08.8	29.79	−87.115	185.7
AGUILA1	October 12, 1979	06:01.8	22.41	−159.71	334.2
AGUILA2	October 12, 1979	07:11.5	22.58	−159.71	442.76
AGUILA3	October 12, 1979	09:56.8	22.97	−159.702	536.36
GAIL	December 4, 1980	07:35.8	29.32	−87.42	181.3
HOPE	December 6, 1980	07:37.9	29.26	−87.041	182.6
IRIS	December 8, 1980	13:07.3	28.76	−87.185	182.2
JAN	December 12, 1980	13:41.6	29.18	−86.978	183.7
LADY LOU	March 29, 1981	57:03.7			180
KLONDIKE KATE	April 3, 1981	6:24:58			180
COLORED BUBBLES	September 17, 1982				
COLORED BUBBLES	September 18, 1982				
CRRES AA-3 A	May 25, 1992	19:52	18.97	−66.6	250

THE MONEY

Of course, we didn't get to conduct all these fun experiments without money. The Department of Defense and Atomic Energy Commission funded the very earliest Los Alamos high-altitude nuclear tests. Then for many years a provision of the 1963 Limited Test Ban Treaty provided funds. As part of Safeguard C of 1963 Limited Test Ban Treaty, the AEC and its successors maintained ships, labs, rockets, aircraft, and a “dedicated staff” to enable the Government to resume testing nuclear weapons in the atmosphere. From a professional point-of-view this funding had the disadvantage that the goal was not to carefully document

the results of the experiments. We had to scramble to develop experiment plans for the next exercise.

Over the years Los Alamos has also received funding from NASA and the Defense Nuclear Agency. Our experience with active experiments resulted in Los Alamos being funded by the Strategic Defense Imitative Office to conduct the B.E.A.R. project in the 80's.

CATALOG

We used the word “catalog” to describe our involvement in active experiments; this report will not describe all of them. They begin

TABLE 4 | Dates and locations of LANL-involved shaped-charge barium injections.

Event	Date	Time (UT)	Latitude	Longitude	Altitude (km)
ALCO	October 18, 1971	15:14:47	22.86	−160.02	466
BUBIA	October 19, 1971	15:15:10	23.10	−160.15	468
OOSIK	March 3, 1972	6:58:59	66.47	−147.52	544.15
CHACHALACA	October 9, 1972	12:58:01	66.23	−148.03	467
LORO	October 18, 1972	15:06:36	23.13	−160.145	555.3
SKYLAB1	November 27, 1973	15:13:001	66.65	−147.5	561
SKYLAB2	December 4, 1973	15:27:00	66.00	−146.6	558
ISPIDA	March 16, 1974	12:35:00	66.42	−147.081	575.002
TORDO UNO	January 6, 1975	23:56	71.50	−127.1	540
TORDO DOS	January 11, 1975	0:32	72.70	−119.5	507.6
LOXIA	May 14, 1975	14:40:00	23.86	−159.893	360.5
PERIQUITO UNO	November 25, 1975	46:00.5	72.17	−115.667	475
PERIQUITO DOS	November 28, 1975	46:00.5	72.58	−116.23	472.91
BUARO	June 7, 1976	30:01.3	23.33	−159.948	451.8
AVEFRIA UNO	May 8, 1978	11:44	37.62	−116.399	193.36
AVEFRIA DOS	May 18, 1978	11:35	37.70	−116.333	190.35
35.007 GE I	March 30, 1984	11:05:40			
35.007 GE II	March 30, 1984	11:11:25			
35.008 GE I	April 1, 1984	7:24:40			
35.008 GE II	April 1, 1984	7:30:25			

with the Teak nuclear test in 1958 and end with the NASA-sponsored CRRES experiments in 1992. Los Alamos has been involved with 107 active experiments in space, not including any RF modification experiments.

Table 1 shows the dates and locations of LANL-involved nuclear tests in space. Los Alamos involvement began with the “TEAK” nuclear test on August 1, 1958 when the author was in high school. **Table 2** shows the dates and locations of LANL-involved thermite releases from orbit. “Thermite” releases will be described in the section on techniques. **Table 3** shows the dates and locations of LANL-involved thermite releases from rockets. **Table 4** shows the dates and locations of LANL-involved shaped-charge barium injections.

The tables show that we conducted experiments over a wide range of altitudes. This “catalog” does not have the experiment altitude for a number of experiments. The altitude of the experiments ranged from 43 km for the ORANGE test to 33,553 km (over 5 Re) for the CRRES G-8 release.

The tables also show that we conducted experiments over a wide range of latitudes and longitudes. They range in longitude from Johnston Island in the Pacific to the Argus nuclear tests in the South Atlantic. Our experiments range in latitude from the South Atlantic Argus tests to Cape Parry in Canada’s Northwest Territories. Los Alamos has been everywhere!

TECHNIQUES

Los Alamos has employed a wide variety of techniques to conduct active experiments in space. These include explosions such as the Argus nuclear tests and the Waterhole experiments use of

high explosives to inject molecules into the F-region ionosphere. High explosives are more efficient than just dumping liquid water. Los Alamos pioneered the use of shaped charges to vaporize and inject barium vapor. The charges had nickel-lined barium cones and generated barium jets with speeds up to 14 km/s. Los Alamos also conducted the granddaddy of all particle injections with the 1-MeV neutral particle B.E.A.R. beam. Los Alamos also conducted the more traditional thermite and sulfa-hexa-fluoride release experiments.

Nuclear Tests

It would be dismissive of the courage and dedication of those involved to write that the technique for the Los Alamos-related nuclear tests involved putting a nuclear device atop a Redstone or Thor missile and detonating the device at a preset time after launch. Both the missiles and the devices were experimental. For example, the Bluegill test was actually Bluegill triple prime because of missile malfunctions on the first three attempts.

Thermite Barium Releases From Sounding Rockets

Barium vapor released from thermite canisters was the most common active experiment technique. The barium vapor was produced by the exothermic reaction of a pressed mixture of barium metal chips and cupric oxide powder. The normal mixture ratio was 2.5 moles of barium per mole of cupric oxide with an addition of 1.8 percent by weight of barium azide (see “*Chemical Releases from Space Shuttle Payloads*,” Thiokol, Wasatch Division, Ogden, Utah, NAS 5-24052, May 1975). Titanium-boride thermites have also been used. Neutral barium atoms evaporate from the hot barium vapor droplets.



FIGURE 1 | NASA's CRRES G-9 release—amateur's photograph. The CRRES satellite's path is from upper right to lower left. The bright green image is resonantly scattered neutral barium atoms. The ionized barium is the blue streak extending upwards and to the left to the canister release point.

The neutral barium expands as a shell centered on the velocity of the release canister. The shell expands with a speed of order 1 km/sec and a thickness of order 0.25 km/sec [3]. Barium was chosen because in sunlight the neutral barium atom's photoionization time constant is roughly 20 s and both neutral barium atoms and ionized barium resonantly scatter sunlight allowing for remote optical sensing of both species [see [4]]. When the canisters are released from sounding rockets at thermospheric altitudes with essentially no directed velocity the resultant barium ion densities can be quite high. Gonzales [5] reported ion densities approaching $10^7/\text{cm}^3$ for hours after the HOPE release at 182 km). This technique is limited because the releases must occur in twilight when the canister is in sunlight and the optical observations are in darkness.

Shaped-Charge Injections

From a basic physics point-of-view the sounding rocket release were excellent for mapping ionospheric winds and electric fields, but they were inadequate for tracing magnetic field lines above the thermosphere. With the encouragement of Gene Wescott from the University of Alaska's Geophysics Institute, Los Alamos pioneered the use of shaped charges to vaporize and inject

barium vapor. The charges had nickel-lined barium cones and generated barium jets with speeds up to 14 km/s and the barium ions could be observed well into the magnetosphere. Another useful feature of the shaped charge injections was that not all the barium was accelerated and a "stay behind" cloud could be tracked to test for altitude effects on field line convection (and equipotentiality). The Alco, Bubia, and Loro experiments were LANL's first use of barium shaped-charges [6].

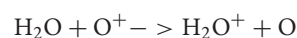
Thermite Barium Releases From Orbit

When we learned that Jim Heppner from Goddard Space Flight Center planned the CAMEO (Chemically Active Materials Ejected from Orbit) releases we suggested that we could track the barium from the lower forty-eight. We pointed out that barium released at orbital velocity would have sufficient perpendicular (to B) velocity that the magnetic mirror force would accelerate the barium ions upwards along the geomagnetic field. The physics underlying this technique is well-documented in Heppner et al. [7].

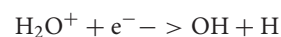
Ammonium Nitrate/Nitro-Methane Explosions

Following Mike Mendillo's paper on the effects of the Skylab launch in May, 1973 [8] we decided to try to duplicate the chemistry attributed to have caused the ionospheric hole. The observations were interpreted in terms of exceptionally enhanced chemical loss rates due to the molecular hydrogen and water vapor contained in the Saturn second-stage exhaust plume.

The F-region ionosphere is dominated by atomic ions (mostly O^+). When molecules are added to the mix there are rapid charge-exchange/dissociative recombination reactions that remove the ions and electrons producing an ionospheric "hole."



and then



We began by considering ways to inject water/steam from sounding rocket, even to considering launching hot water heaters. However, a colleague, John Zinn, who had a back ground in explosives pointed out that most explosives produce water and carbon dioxide. So, *voila* detonate a high explosive in the F-region ionosphere!

After some study of products of detonation from several explosive mixtures considering safety and maximizing the production of molecules, we settled on ammonium nitrate ("fertilizer") and nitromethane ("nitro" for drag racers) for our test. Fortunately, LANL had experts in producing explosives and they pressed ammonium nitrate into a cylindrical tube. They also calculated the correct stoichiometric mixture of liquid nitro methane to add to the tube on site.

Our initial experiments loaded the ammonium nitrate into aluminum tubes, but concerns about the dangers of falling metal fragments forced us to use plexiglass tubes for later experiments. This caused some consternation for the

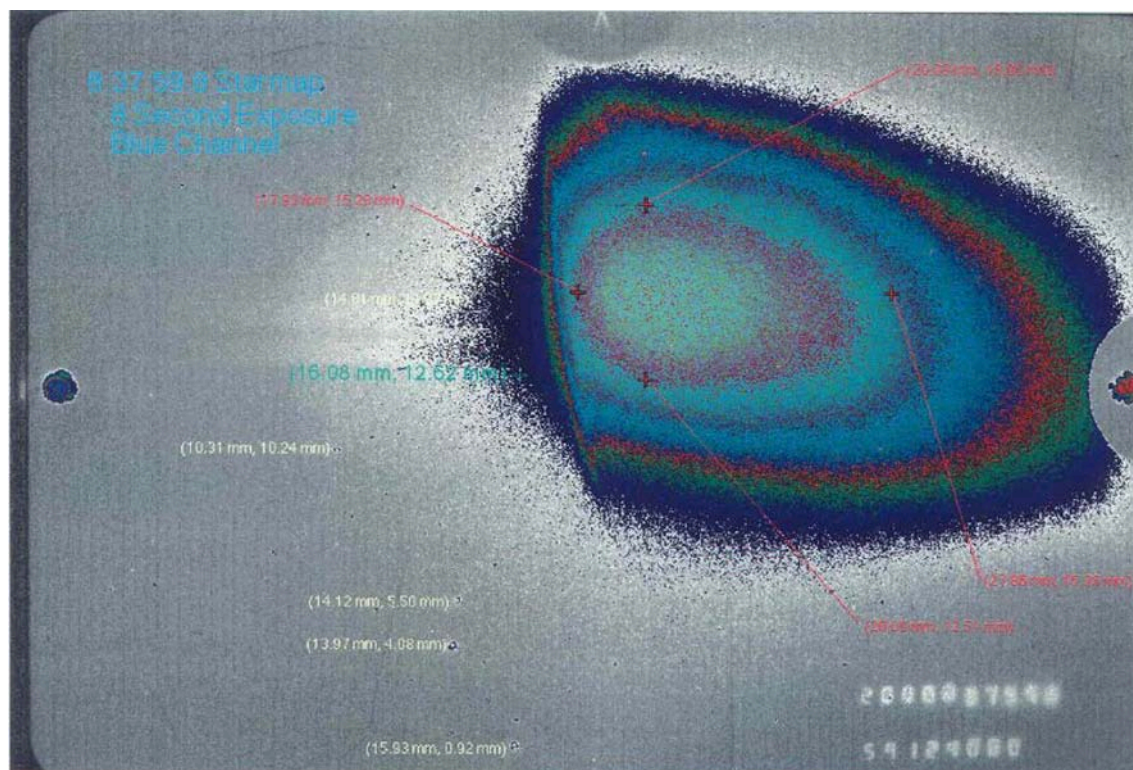


FIGURE 2 | NASA's CRRES G-9 release—Ball (ionized barium) in false color contours. In this view the release satellite had been traveling from left to right. The geomagnetic field appears essentially vertical.

final Waterhole experiment at Canada's Churchill Research Range in Manitoba. We were able to ship the ammonium nitrate tube and the nitromethane ("cleaning fluid") separately to be assembled in an underground bunker at the range. Unfortunately, as we were topping off the appropriate amount of nitromethane into the tube it began to leak. After much perspiration and yellow ("rocket tape") we were able to finish the assembly.

Particle Accelerator

The first particle accelerator flown by Los Alamos was during Operation Birdseed in 1970. The accelerator was a co-axial, neon plasma gun designed by John Marshall and Ivars Henins. The energy delivered to the plasma gun was 350,000 joules at 1,700,000 amperes and a power of 40 billion watts [9]. A number of active experiments in space have employed electron guns, but to my knowledge Los Alamos was not involved in the employment of those devices.

Los Alamos' next venture into the particle accelerator technique came in the 80's at the request of President Reagan's Strategic Defensive Initiative (SDI). Neutral particle beam (NPB) technology was considered to be one of the most promising SDI concepts. The challenge was to fly a Radio Frequency Quadrupole (RFQ) accelerator on board a sounding rocket. The accelerator first accelerated negative hydrogen ions to 1 MeV and then passed them through a gas to strip off the electron resulting

in a 1-MeV NPB that would propagate across the geomagnetic field [10].

DIAGNOSTICS

Diagnostics are the key to successful active experiments. Over the years optical diagnostics have been the backbone of our active experiments. We have employed both ground-based and airborne platforms. The advantage of the airborne platforms is the cloud-free lines of sight to the experiment. The so-called "Readiness to Test" program funded Boeing 707s for airborne optical diagnostics at all three AEC laboratories. Los Alamos has employed *in situ* diagnostics since the nuclear testing days; the Argus experiments were diagnosed by instruments on Explorer IV. We got back to *in situ* measurements during the 1976 Buaro shaped charge experiment when Harry Koons measured the electric fields generated by the free energy of the barium ions [11]. Of course, the satellite-borne sensors made crucial diagnostics of the CRRES releases. The original ionospheric depletion experiments, Lagopedo Uno and Dos, in 1977 were measured by ionosondes located on the island of Kauai. Some of the CRRES experiments employed RF diagnostics of the Arecibo facility.

Cameras provided the principal diagnostic for our active experiments. Television cameras provided real-time tracking of barium clouds out to several earth radii distance. Of course, film

CRRES BARIUM RELEASE G-12 08-12-91 09:31:18 UT

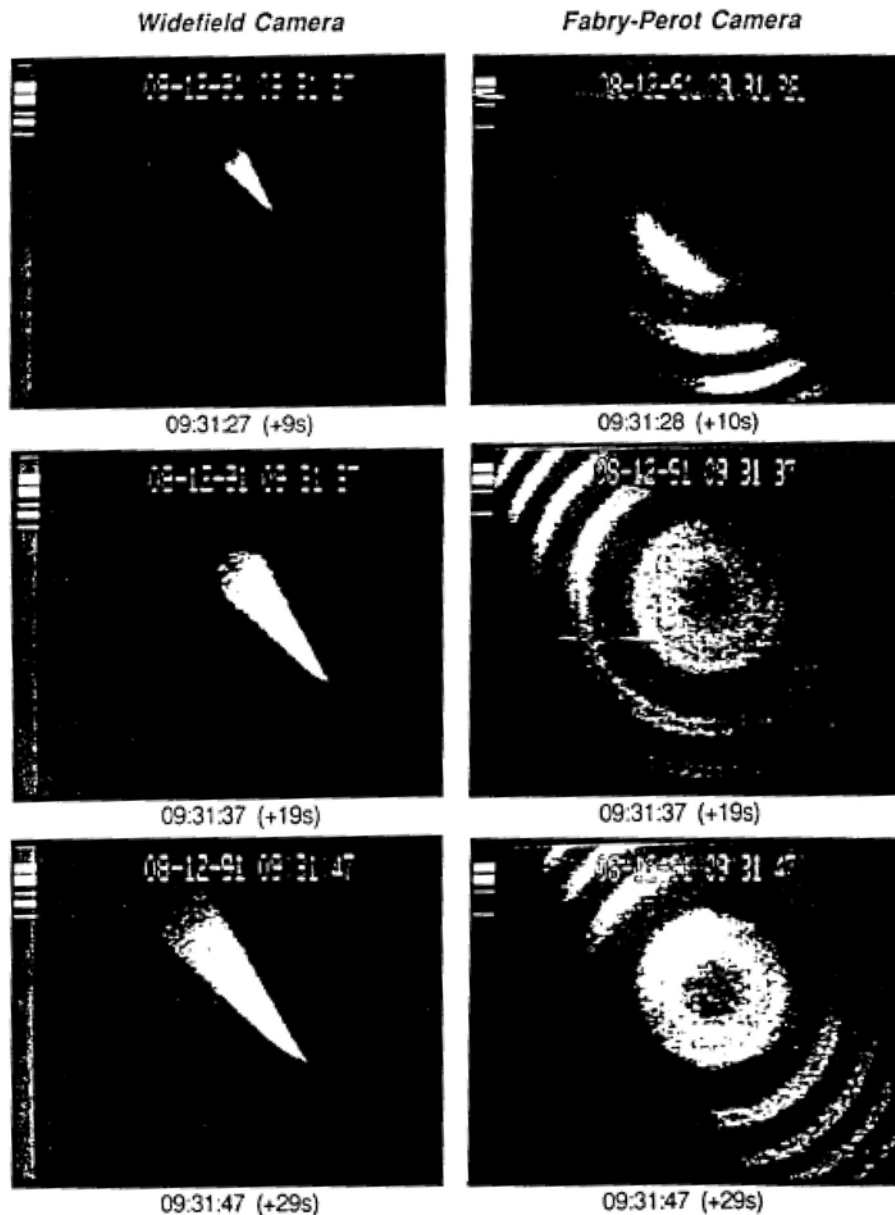


FIGURE 3 | G-12 barium release filtered image sequence. Fringes from the 2-mm Fabry-Perot etalon...reveal the double-peaked nature of the ion radial velocity distribution [14].

cameras provided quantitative image data needed for inventories and dimensions. All-sky cameras provided back-ups. Because distant images are faint and we need star background for triangulation we used interference filtered image intensifiers for the barium-related experiments. Rick Rairden's airborne Fabry-Perot allowed us to sense barium ion motions (not just locations) and provided unique confirmation of barium ion magnetization (ions moving toward and away from the sensor). We also fielded spectrographs and photometers.

As the saying goes “A picture is worth a thousand words—and takes gigabits to process!” [see [12]]. Even amateur photographs can provide valuable information. **Figure 1** was taken by my daughter on the beach in St. Croix. It shows the dramatic G-9 CRRES release. The neutral barium atoms are imaged by the bright green sphere. The trailing blue light comes from ionized barium. The cloud in this photo also demonstrates the limitations of ground-based optical diagnostics. **Figure 2** shows a quantitative measure on ionized barium column density. This

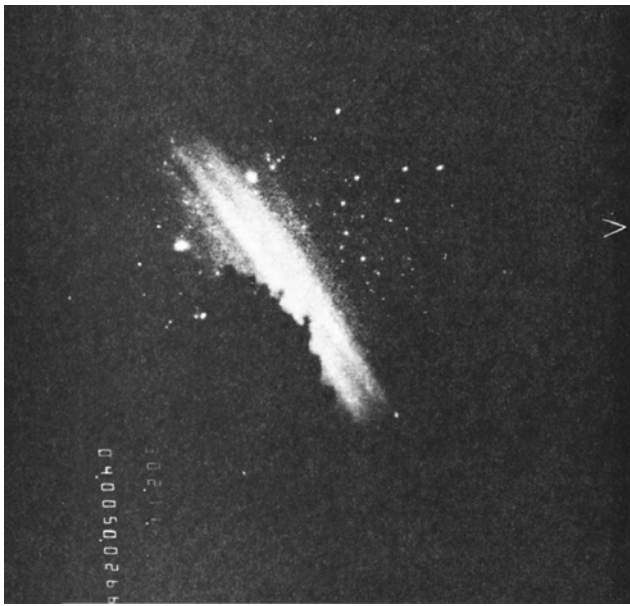


FIGURE 4 | Intensified camera view of release 4 over trees (left side) from Table Mountain, California, at 1142 UT. Table Mountain Observatory filtered, intensified camera image of the CAMEO polar cap barium release showing star field and obstruction by tree [Heppner, et al. [7]; reprinted with permission of American Geophysical Union].

image also shows the slight curvature to the back-side of the barium cloud as explained by Delamere et al. [13].

Figures 3, 4 show two examples of the specialized optical diagnostics we employed. **Figure 3**, from Rairden et al. [14], shows the G-12 barium ion images at three times. On the right we see images from the co-aligned Fabry-Perot instrument. We want to emphasize the middle donut-shaped image. The displacement from the donut hole is a measure of barium ion velocity. The dimple is the first fringe reveals “the double-peaked nature of the ion radial velocity distribution.” Voila—magnetized barium ions!

The intensified camera image in **Figure 4** shows one of the field-aligned CAMEO barium streaks. This image was captured with an interference-filtered intensified camera located at Table Mountain Observatory near Los Angeles, California. The thermite barium release from a satellite occurred over the North Slope of Alaska and the magnetic mirror force pushed the barium ions up the field line to where they were detected thousands of kilometers away. This image also shows another disadvantage of ground-based photography—note the tree obstructing part of the barium streak. In fact, the folks at Table Mountain had their instruments located in the back of a U-Haul truck and had to shove the trackers further toward the back of the truck when the tree became a problem. This demonstrates a unique challenge faced by those diagnosing active experiments—the skill and resources of the experimenter play a role in real-time data acquisition.

We also employed computer modeling of the images. They were helpful in experiment planning to determine camera

pointing, field-of-view, and brightness. Computer modeling was also essential in understanding the phenomena being measured.

OBJECTIVES

The nuclear tests in space were instrumental in testing device designs, studying weapon effects, and testing delivery systems. The weapon effects included enhanced ionization, diamagnetic cavity formation and collapse, electro-magnetic pulse generation, electro-magnetic wave propagation, atmospheric heave, energetic particle motions and trapping. They were considered for ICBM defense as well as radar blackout studies.

The earliest barium release experiments were used to measure convection electric fields and winds in the ionosphere. Next the shaped-charge experiments were used to illuminate high altitude magnetic field lines and their convection. Injections of energetic barium ions confirmed the magnetic mirror force on ion dynamics.

Then we got more adventurous in our experiment objectives. Active experiments provide unique opportunities to study fluid and kinetic plasma instabilities. Most plasmas encountered in nature are close to equilibrium and not likely to be unstable. With active experiments we can “set the ball at the top of the hill” and watch it fall down. The images in **Figure 5** show the “up-the-field-line” images of the Avefria Dos barium cloud. This experiment occurred in Nevada allowing us to position cameras at the foot of the field line. The images show prompt structuring of the energetic barium plasma jet on the left and the slower, Rayleigh-Taylor fluid instability structuring of the “stay behind” barium ions on the right.

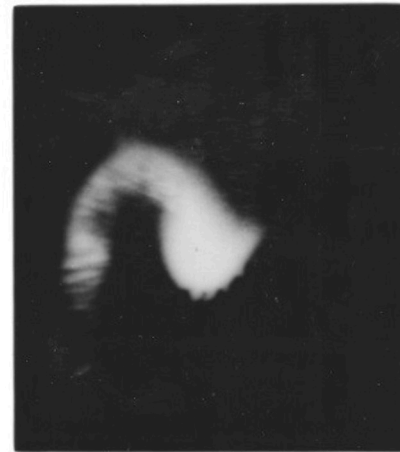
Figure 6 shows the spectrogram of the field intensities for the G-9 chemical release on July 19, 1991. On the top is the spectrum of the magnetic field fluctuations and on the bottom half we see the spectrum of the electric field fluctuations. This data come from sensors on the CRRES satellite flying through the barium cloud moments after release; the broad-band intensification shown in pink marks the event. The data is from Koons and Roeder [16]. At one time we claimed that our active experiments would make the space plasma “ring like a bell”; the data show predominantly broad-band electrostatic emissions and not bell-like resonant tones.

Our active experiment objectives included the study of many additional phenomena. These included Critical Ionization Velocity (CIV) studies—a hypothesis predicted by Hannes Alfvén to account for the composition of solar system planets. The critical ionization velocity for a neutral cloud to become ionized is when the relative kinetic energy is equal to the ionization energy. Another objective was to test models of RF propagation through structured plasmas—the PLACES experiments. We also conducted experiments to test the relationship between thermal electron currents and auroral electron precipitation—the Waterhole experiments. We studied the formation of diamagnetic cavities and polarization electric fields in the CRRES experiments. Many of the barium experiments were used to simulate High Altitude Nuclear Explosion (H.A.N.E.) phenomena.

"RAYLEIGH-TAYLOR" INSTABILITY



a. Release +2.0 min.



b. Release +2.5 min.



c. Release +3.0 min.



d. Release +3.5 min.

"UP-THE-FIELD-LINE" VIEW OF AVEFRIA DOS BARIUM EXPERIMENT (Looking upwards - Magnetic Zenith).

FIGURE 5 | Images of structuring in barium ion clouds. These sequenced "up-the-field line" images show development of a Rayleigh-Taylor instability on the right. The barium ions imaged on the left came from the energetic barium injected by a shaped-charge. That structuring was probably due to a kinetic instability [15].

EXAMPLES

This review cannot possibly describe all 107 Los Alamos active experiments so we'll use examples to describe the breadth of our work. We begin with the Orange Nuclear Test. Then we show data from the field-line tracing experiments Tordo and Periquito. Following that we'll cover a unique series of the experiments—the Waterhole ionospheric depletion experiments. Then, we'll describe barium releases designed to study RF propagation through structured plasmas and we'll review thermite barium releases at orbital velocity, the CAMEO and CRRES experiments.

We'll close with the most energetic particle the B.E.A.R. 10-mA (equivalent), 1-MeV, neutral hydrogen beam.

Orange—Nuclear Weapons Effects Test

The Orange, nuclear weapons effects test, was conducted on August 12, 1958 as part of Operation Hardtack I [17]. The 3.8 megaton device was exploded 43 kilometers above Johnston Island in the Pacific. **Figure 7** shows the Orange Event at 1 min after detonation. Note the toroidal yellow or orange colored fireball and white- blue-green-purple air radiation induced glow. This photograph was taken from the deck of an aircraft carrier.

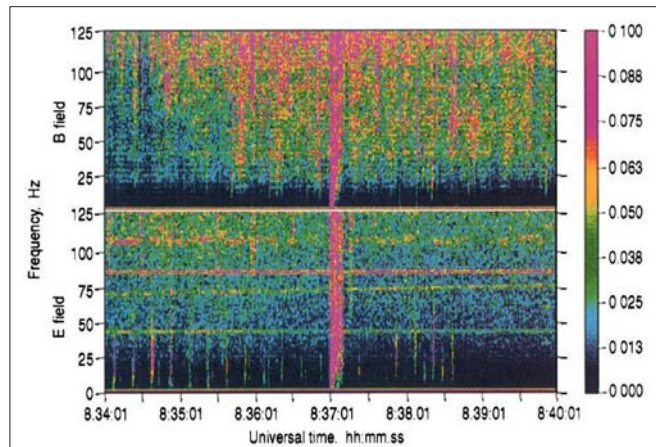


FIGURE 6 | Spectrogram of the field intensities for the G-9 chemical release on July 19, 1991. **(Top)** Magnetic field; **(Bottom)** electric field. [[16]; reprinted with permission of American Geophysical Union].



FIGURE 7 | Orange event: toroidal yellow or orange colored fireball and white-blue-green-purple air radiation induced glow photographed from the deck of a U.S. aircraft carrier at 1 min after burst, 12 August 1958 [17].

One might question citing this as an example of an experiment in “space,” but, in fact, the large energy release caused “heave,” an upwelling of the neutral atmosphere into the thermosphere. We have heard of some exotic techniques proposed to “dump” anomalous levels of satellite killing radiation. We suggest that the neutral atmospheric “heave” from a low altitude, high yield explosion would “heave” a massive quantity of neutrals into the upper atmosphere causing energetic particles to scatter

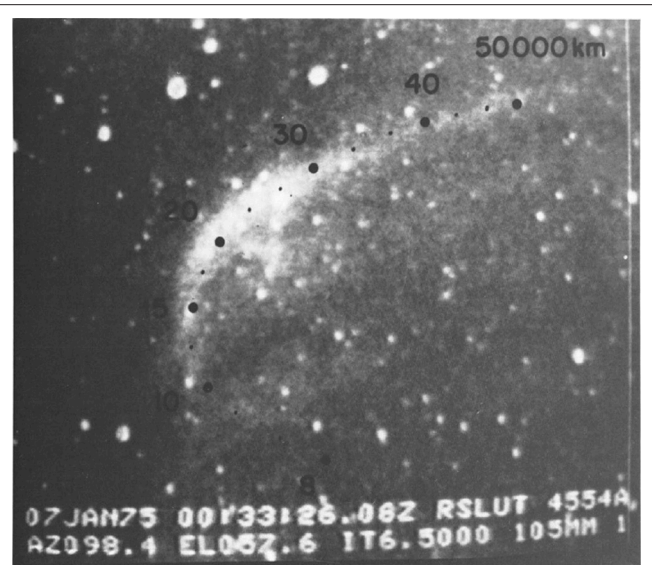


FIGURE 8 | TV frame and theoretical field line from resolute Bay, Northwest Territories at 0033:26 UT. The normal termination at the upper end of the streak (lower right) would near 27,000 km. Note that the apparent end is near 50,000 km. Barium illuminated field line extending to 8 Re altitude. [[18]; reprinted with permission of American Geophysical Union].

and precipitate. Detonation at a location conjugate to South American Anomaly would result in explosion-produced betas being quickly dumped.

Field Line Tracing 1975 Shaped Charge Injections From Cape Parry Canada—Tordo and Periquito

Figure 8 describes how shaped-charge barium injections were used for field-line tracing. In January and then again in November of 1975 Los Alamos working with our Sandia, EG&G, Canadian and Geophysical Institute partners launched rockets from Cape Parry, Northwest Territories. Shaped-charges carried aboard these rockets injected barium ions up the field lines into the polar cusp region. The TV image (**Figure 8**) from Wescott et al. [18] shows a streak of barium ions extending thousands of kilometers (about 8 Re) into space. **Figure 9** from Jeffries et al. [19] shows the track of the leading tip of ionized barium streak for the Tordo Uno injection, projected down along magnetic field lines to the 100 km reference altitude. Numbers along the track are minutes after injection. Note the clear demonstration of anti-sunward convection over the polar cap. Dungey was right!

Field Line Tracing CAMEO

The next example of field-line tracing is the CAMEO (Chemically Active Materials Ejected from Orbit) experiment of Jim Heppner. Jim told us that he had arranged for thermite barium releases over Alaska from a polar orbiting satellite. We replied that we’d track the barium ions from the lower 48 relying on the magnetic mirror force on the energetic barium ions to overcome gravity and lift the ions upwards along the magnetic field. Indeed, barium streaks photographed from Mount Haleakala, Hawaii

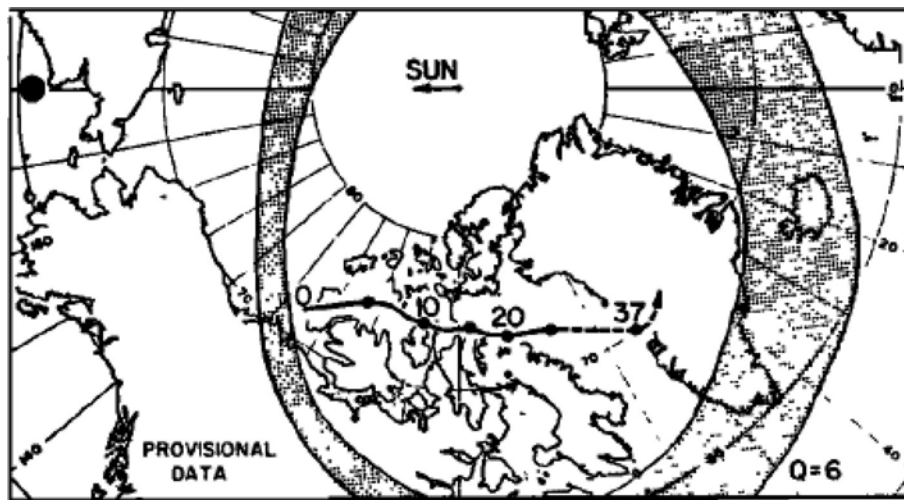


FIGURE 9 | Track of leading tip of ionized barium streak for Tordo I, projected down along magnetic field line to 100 km altitude. Numbers along track are minutes after injection; circles are at 5-min intervals except for 37-min location Dashed line after 25 min reflects present uncertainties requiring additional analysis The statistical auroral oval for a distributed $Q=6$ magnetic index is shown as the shaded area. Anti-sunward convection over the polar cap [19]; reprinted with permission of American Geophysical Union].

and Table Mountain Observatory in California were triangulated measuring the ion motion upwards along the magnetic field line. **Figure 10** shows the altitude of the release number two ions as a function of time. Accelerations parallel to B were required to account for the barium ion position as a function of time. In fact, the trajectory indicates up to 6 keV $E_{||}$ acceleration and deceleration.

Plasma Depletion Experiments—Waterhole I and III

The Waterhole experiments were ammonium nitrate/nitro methane explosions in the auroral F-region. Charge-exchange/dissociative recombination chemistry removes ions and electrons forming a 50-km diameter “hole” in the ionosphere. The hypothesis was that field-aligned currents connected to auroral arcs are important to the mechanism producing the arc and removing the thermal plasma will perturb the currents and modify the acceleration mechanism [20].

The Waterhole experiments utilized what we learned about depleting the ionosphere following the Skylab launch and Los Alamos’ Lagopedo experiments. Releasing tri-atomic molecules in the O^+ dominate ionosphere leads to rapid charge-exchange and then dissociative recombination chemistry which removes ions and electrons forming a 50-km diameter “hole” in the ionosphere. It turns out that water (H_2O) is a great molecule to release. It turns out that a nitro-methane/ammonium nitrate (basically fuel oil and fertilizer) mixture also works great. You pack the ammonium nitrate into a cylinder and then add the liquid nitro-methane to get the correct stoichiometric balance.

So, in April of 1980 we flew an 88-kilogram ammonium nitrate/nitro-methane explosive into the aurora above Churchill, Canada. **Figure 11** shows *in situ* data obtained by our Canadian partner Whalen et al. [22]. Curve (a) shows rocket altitude and

distance from event, curve (b) shows the relative local electron density with a dramatic reduction until the payload flies out of the hole, curve (c) shows the precipitating electron intensity at 0.5 keV again with a dramatic reduction until the payload flies out of the hole, and curve (d) shows the peak column emission intensities of auroral green line. Our question was, “Did we turn off the aurora?”

So, with the interesting Waterhole I results we were able to secure funding the try again. By the way Waterhole II suffered a rocket malfunction and the high explosive landed a few kilometers away from our Churchill ground station where the Mounties detonated it with a shaped charge. It turns out that on Waterhole I we detonated the high explosive just north of the auroral field line. For Waterhole III we had additional high explosive and the detonation was controlled from the ground when we encountered the precipitating electron flux.

And, of course, we got different results. The precipitating electron flux at 1.5 keV was enhanced at small pitch angles! Quoting Whalen et al. [21], “The rapid response...and spectrum changes...in energetic electron precipitation indicates...induced electric field must have been large enough to accelerate electrons up to several keV” and “Although the two results appear to be contradictory, simple models...of the structure of auroral arcs seem to be in agreement with both experiments.”

These experiments should be repeated perhaps with the launch of liquid fueled rocket passing perpendicular to an auroral arc. This would ensure that thermal electron currents over the arc and on each side were disrupted. An explosive release creates a deep, localized hole in the F-region. The spatially extended release from a rocket burn would be deep enough but more extensive. An experiment in view of Alaskan incoherent scatter radars and ground based all-sky cameras would provide better diagnostics.

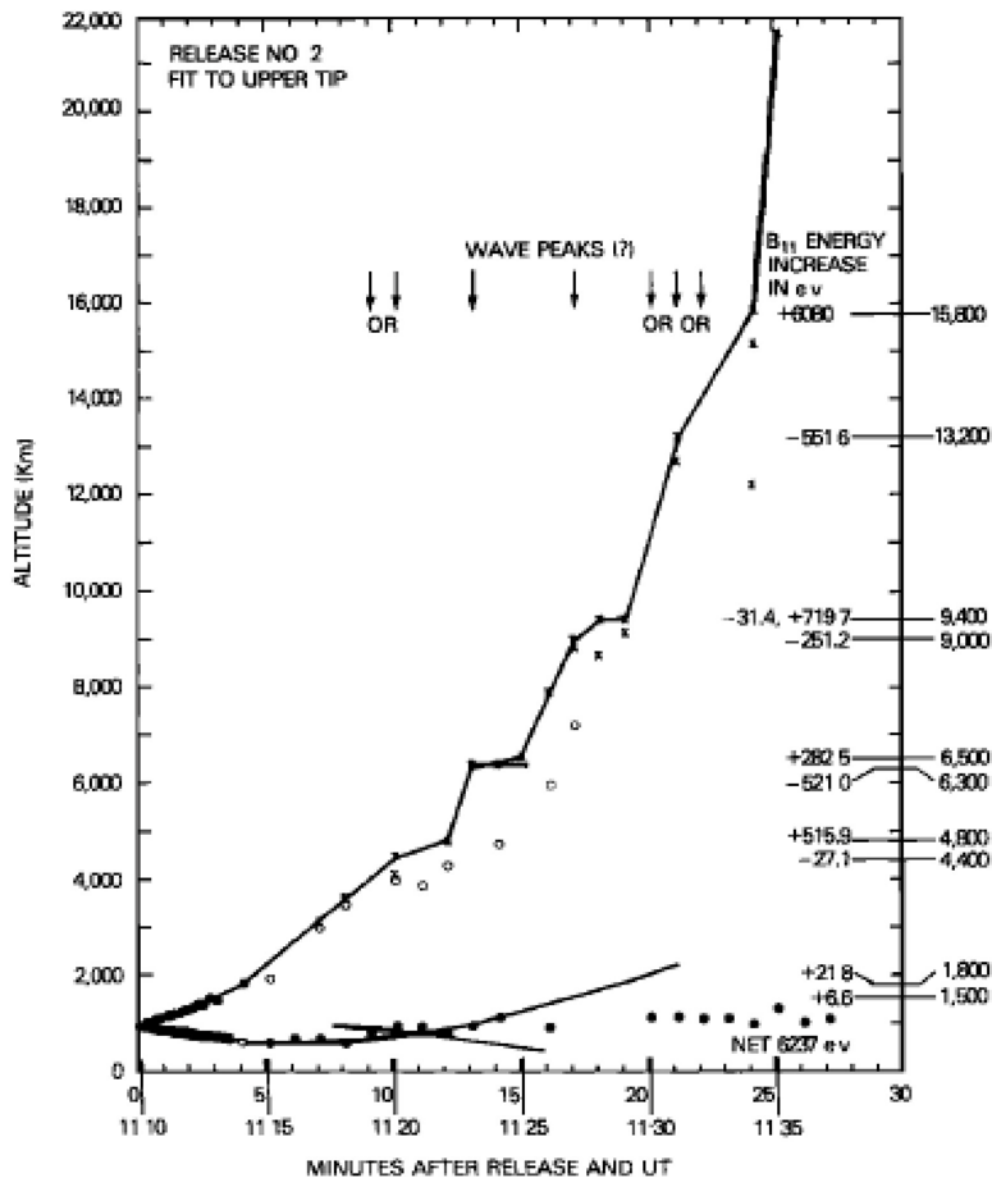


FIGURE 10 | Altitude vs. time plot of CAMEO release No. 2. Trajectory indicates up to 6 keV E|| acceleration and deceleration [7]; reprinted with permission of American Geophysical Union].

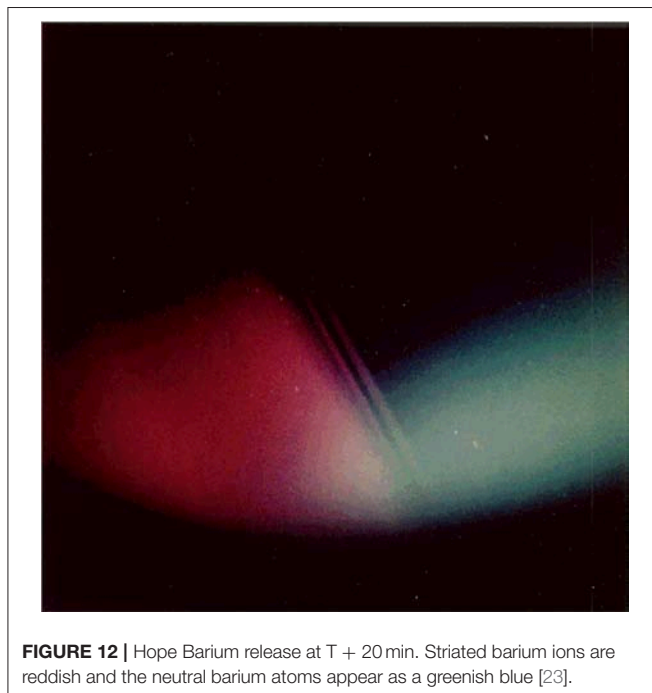
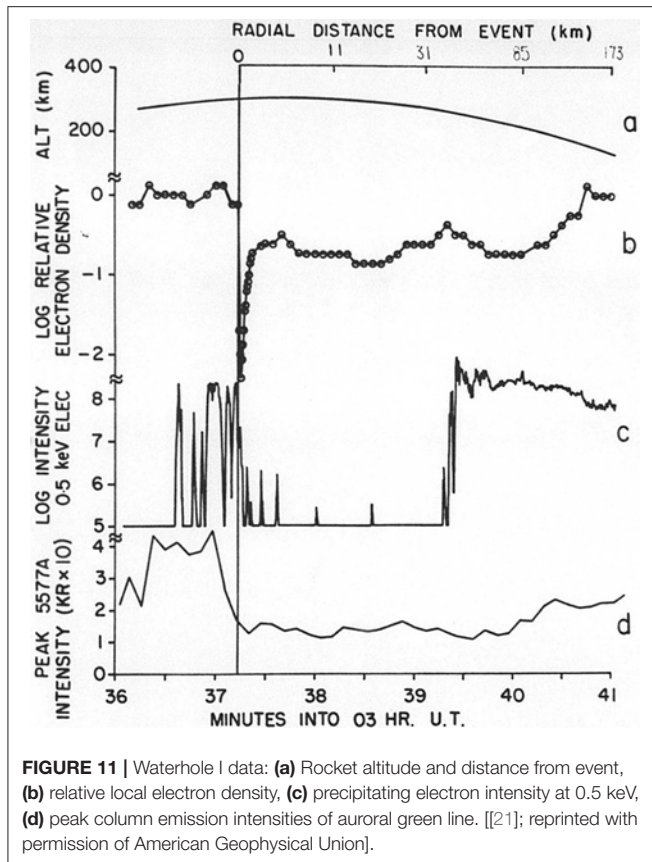
Thermite Barium Releases

Thermite Barium Releases in the Ionosphere

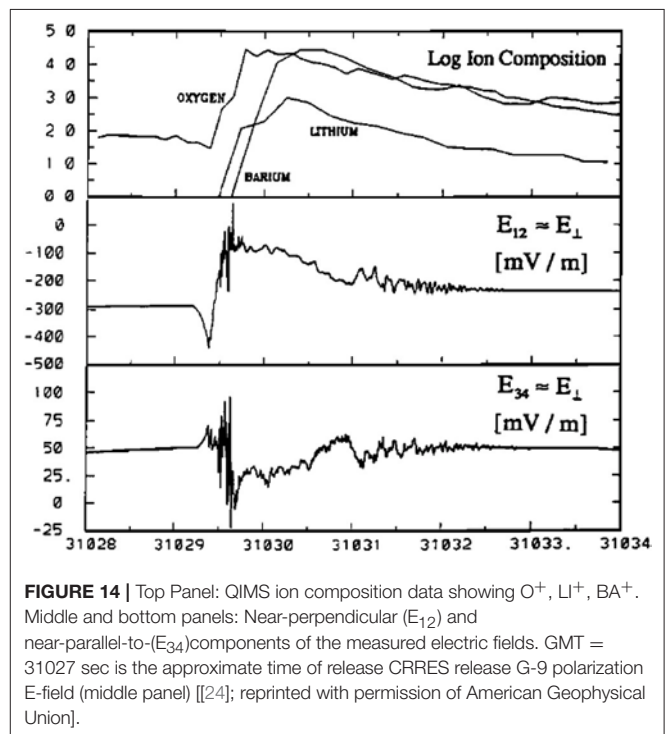
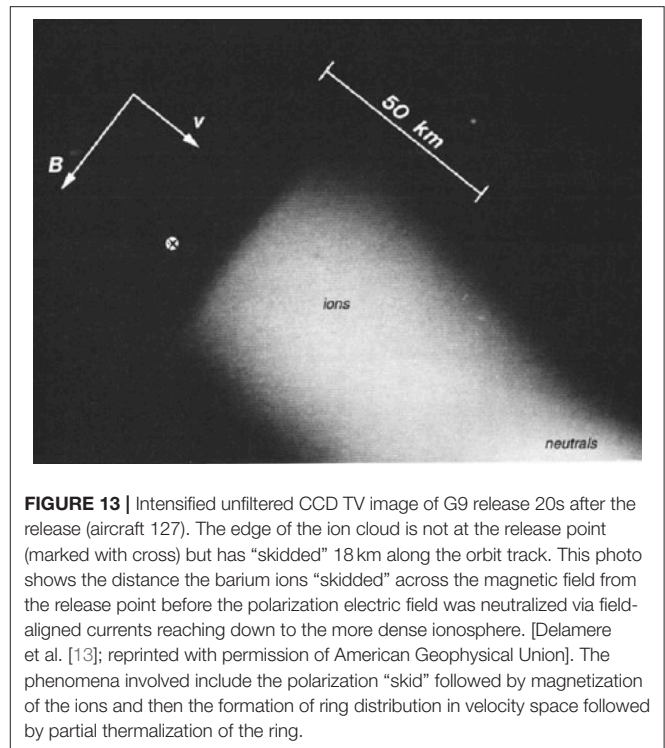
Next, we'll describe our use of thermite barium releases to create a structured plasma. The PLACES (Position Location and Communications Effects Simulations) experiment was a communications field experiment carried out by the Defense Nuclear Agency (DNA) to investigate the effects of structured (striated) ionospheric plasmas on transionospheric communications links (satellite to ground and vice versa). The experiments were carried out in December 1980, at Eglin Air Force Base, Florida. The structured plasma was produced by releasing 48-kg charges of barium thermite near 185-km altitude in the late evening F-region ionosphere on 4 separate days.

The resulting barium plasmas form field aligned structures or striations in the ionosphere that simulate important features of the striations produced by debris plasmas resulting from high-altitude nuclear explosions (see Figure 12).

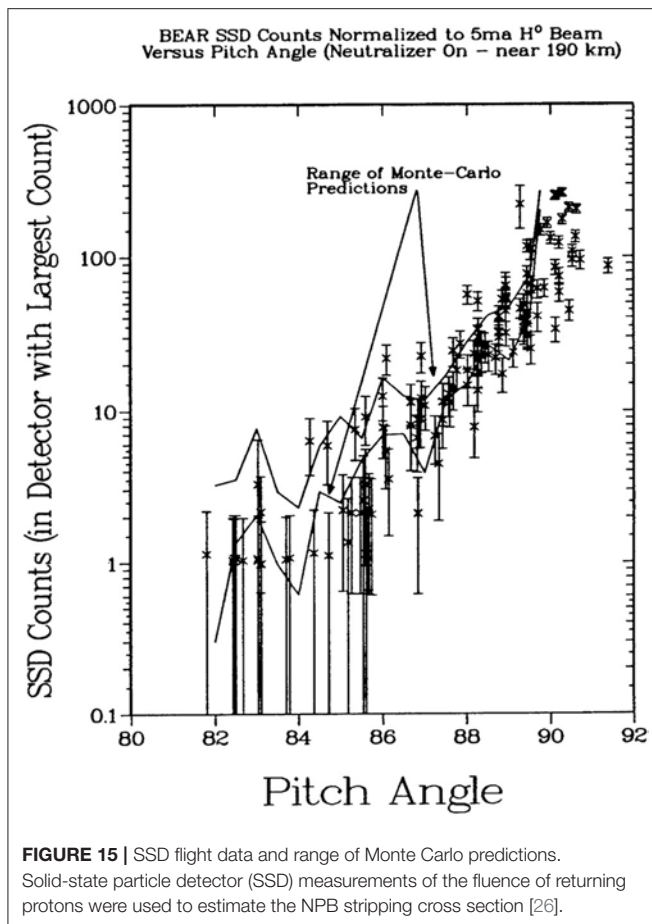
The primary objectives of the PLACES experiments were to determine, by direct measurement, the phase and amplitude scintillations induced by the disturbed plasma upon satellite signals. Simultaneous measurements of the actual plasma structure and spatial distribution by *in situ* and remote diagnostics would then define the true plasma configuration. Extensive theoretical work on the relationship between scintillations and plasma structure would then be open to detailed comparison with these data.



The diagnostics include optical measurements of the time evolution of the power spectral density (PSD) of striations for the electron column density perpendicular to the magnetic field and



measurements of the time-of-arrival spread of energy (channel impulse response) on a phase coded spread spectrum signal emanating from a rocket launched behind the barium cloud and received at specially constructed ground receiving site in



northern Florida (Beacon experiment). The results demonstrated success: the data are shown to be in good agreement with the DNA propagation channel model and a geometric optics interpretation of the observed propagation effects [12, 23].

Thermite Releases at Orbital Velocity

The next active experiment example is the CRRES G-9 experiment, a thermite barium release from the CRRES satellite moving at orbital velocity (about 10 km/s) perpendicular to the local geomagnetic field. This experiment was conducted above the US Virgin Islands on July 19, 1991. A color photograph (Figure 1) showed the dramatic appearance for anyone looking at the right place at the right time. Using the Figure 13 photo Delamere et al. [13] describe the “skid” of the barium ions across the magnetic field. The data on Figure 14 from Szuszcwicz et al. [24] shows the polarization electric field that allowed the barium ions to “skid.” Huba et al. [25] provided a quantitative description of the “skidding.” Rick Rairden’s Fabry-Perot data (Figure 3) showed the ions first skidding, and then becoming magnetized and finally thermalizing. Recall that Figure 6 showed electrostatic field enhancements from the kinetic plasma instability [16]. A smorgasbord of plasma physics!

B.E.A.R.—Particle Beam

We’ll close with a description of our most “active” active experiment, the B.E.A.R. NPB test conducted in July 1989 from White Sands Missile Range. So, we’ve gone from Resolute Bay in January to White Sands in July! The B.E.A.R. experiment was in support of President Ronald Reagan’s Strategic Defense Initiative (SDI). The challenge was to fly a RFQ accelerator on board a sounding rocket. The accelerator first accelerated negative hydrogen ions to 1 MeV and then passed them through a gas to strip off the electron resulting in a 1-MeV NPB that would propagate across the geomagnetic field. We believe that a 1-MeV beam is the most energetic ever flown in space by about a factor of 30!

Our task was to develop a beam diagnostic package that would measure beam energy, current, divergence, beam composition, beam pointing and beam propagation before stripping. We measured beam pointing well enough to know whether we were aimed at the top or bottom half of the Washington Monument from White Sands. We also monitored spacecraft charging because there was concern that the rocket body would charge up and not allow the beam to propagate away. To do that we alternately turned on and off the neutralizing gas to create a negative ion beam and we also over-neutralized the beam to produce a proton beam. Hugh Christian’s electrostatic analyzer measured spacecraft charging [26].

Measuring how far the neutral hydrogen atoms traveled before suffering a stripping colliding with the atmosphere required a bit of trickery. We had no target to shoot at so we relied on the magnetic mirror force to bring the protons back to Ted Fritz’s solid-state detector on the rocket. We used the rocket ACS to fire the H_0 beam down and east.

Stripping produced protons that mirror and drift up and west back to the rocket. There is a one-to-one relationship between the pitch angle of an observed proton and the distance it traveled as a hydrogen atom before stripping. Figure 15 shows solid-state particle detector (SSD) measurements of the fluence of returning protons and the range of counts predicted by Joe Fitzgerald’s Monte-Carlo code.

The experiment successfully demonstrated that a particle beam would operate and propagate as predicted outside the atmosphere and that there are no unexpected side-effects when firing the beam in space.

EPILOG

Over the years Los Alamos has successfully employed “hypothesis testing” active experiment techniques to complement their satellite-based “discovery” approach to understanding space plasmas. We urge funding agencies to reinvestigate this method of study.

AUTHOR CONTRIBUTIONS

The author confirms being the sole contributor of this work and has approved it for publication.

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Overview of APEX Project Results

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The APEX mother-daughter project (Active Plasma EXperiments) was launched into an elliptical polar orbit (440–3,080 km) in December 1991. It consisted of the main Russian Interkosmos–25 (IK–25) satellite and the Czech MAGION–3 subsatellite, both with international scientific payloads. The mission used intensive modulated electron beam emissions and xenon plasma or neutral releases from the main satellite for studies of dynamic processes in the magnetosphere and upper ionosphere. Its main scientific objectives were to simulate an artificial aurora and to study optical and radio emissions from the aurora region, and to investigate the dynamics and relaxation of modulated electron and plasma jets, artificially injected into the ionospheric plasma. The experiments studied the Critical Ionization Velocity phenomenon and a diamagnetic cavity formation during the xenon releases, local and distant effects of the electron beam injection, spacecraft charging and potential balance, and plasma-wave interactions during the artificial emissions. Attempts were performed to utilize the modulated electron beam as an active transmitting antenna in the space. The theory of ballistic wave propagation across plasma barrier was tested in a joint active experiment with the Dushanbe ionospheric heater facility. In the paper, we give a short overview of the IK–25/MAGION–3 scientific instrumentation and methodology of experiments with artificial beam injections and we provide a review of the main APEX active experiments results, many of which have been published only in the Russian language so far. From a historical 25-years-long perspective, we try to put the results of the APEX experiments into the context of other active experiments in the space plasma.

Keywords: active plasma experiment in space, electron beam, xenon plasma injection, neutral xenon injection, plasma waves, APEX, INTERCOSMOS–25, MAGION–3

1. INTRODUCTION

Active experiments in space plasma utilize many different agents to disturb the ionospheric or magnetospheric environment and to stimulate many variable but rather rare natural phenomena under controlled conditions (Raitt, 1995). The response is then studied in order to get information on natural space structures (e.g., electron beam tracing of magnetospheric or ionospheric electric and magnetic fields), to use artificially produced phenomena as models of natural ones (artificial aurora etc.) or to verify mechanisms which would explain some physical processes (e.g., beam–plasma interactions, wave generation, photochemical reactions) when the relevant physical conditions cannot be achieved in laboratory experiments. Some active experiments have close connections to technology. They are motivated by problems in the design and reliability of the artificial satellites or they help to develop new power generators or propulsion methods.

Since 1960s, tens of active charged particle beam experiments (e.g., Winckler, 1980; Grandal, 1982; Podgorny, 1982; Neubert and Banks, 1992; Raitt, 1995), in space have been performed, the motivations behind these experiments ranging from investigations of charge neutralization processes in the plasma medium near the beam source to probing conditions on remote sections of geomagnetic field lines (Winckler, 1980). A lot of interesting experimental results was obtained aboard the U.S. space shuttles (e.g., Burch, 1986), few other satellites were devoted to active experiments in space plasma at altitudes starting in ionosphere to solar wind. During last two decades their frequency has decreased and the scientific community has turned more to theoretical studies, computer simulations, or data reanalysis in this field. New science mission concepts (e.g., MacDonald et al., 2012) may return active space plasma experiment to the foreground of scientific interest.

The topic of this overview paper—the APEX project satellite pair—had relatively unique orbit among the active experiment spacecraft as it will be described later. It had extensive scientific payload and ambitious science plans. Although it is more than 26 years since the launch now, new original research articles still profit from the collected data. We bring an overview of the active APEX experiments and their results, while many other interesting results of passive ionospheric observations (often based on the two-point satellite measurements rare in that era) are out of the scope of this paper. The overview should not be considered as a complete list of references to the APEX project related literature, some results were originally published in Russian and only later appeared in a Russian translation version of the journal or in other English language journals or proceedings and for such cases we usually cite just one reference.

The review is organized to several sections. The first one comprises introduction to the APEX project, its origin, scientific goals, and a short summary of the scientific instruments onboard the two satellites of the project. Comments to the active experiment planning and telemetry issues are included here as well. The result overview section starts with analysis of the main satellite charging during the passive observations and active experiments with the electron beam injection neutralized by neutral xenon/xenon plasma releases or during the xenon plasma releases itself. Overview of the neutral xenon releases related to the critical ionization velocity (CIV) phenomenon processes follows. Next part is devoted the beam-plasma interaction, discussing the plasma environment in a vicinity of the main satellite during the beam injections, diamagnetic cavity formation, electron pitch-angle distributions, and generation of ELF/VLF and HF waves. Distant injection effects as observed from the subsatellite at distances of hundreds of kilometers are summarized in the next subsection. The last part briefly covers the ionospheric heating experiments in which the main APEX satellite participated. Summary and concluding remarks close the article.

2. APEX PROJECT

The international experiment APEX in frame of the former INTERCOSMOS science program was suggested by IZMIRAN—the Pushkov Institute of Terrestrial Magnetism, Ionosphere,

and Radio Wave Propagation of the Russian Academy of Sciences, Troitsk, Russia, and the Geophysical Institute of Czechoslovak Academy of Sciences, Prague, Czech Republic (GFI, the group is affiliated to the Institute of Atmospheric Physics of Czech Academy of Sciences today) in mid 1980's. The project followed previous Soviet/international active experiments with sounding rockets and satellites (Stereotop, PORCUPINE, Trigger, Zarnitsa, G-60-S, ARAKS, ACTIVE) (Cambou et al., 1975, 1980; Dokukin et al., 1981; Haerendel and Sagdeev, 1981; Sagdeev et al., 1981; Managadze et al., 1988; Klos et al., 1998), this time in a new mother-daughter configuration and with wide international payload (Russia, Ukraine, Czechoslovakia, Poland, Rumania, Bulgaria, Hungary, eastern Germany). The name of the Interkosmos programme project APEX is abbreviated from words “Active Plasma EXperiments.” The core of the project were active experiments with injection of electron and (or) plasma beams from a board of the low Earth's orbit (LEO) satellite with simultaneous registration of the phenomena, produced by interaction of injected beams with background plasma.

Essentially important peculiarity of the project were synchronous measurements of the basic physical parameters of the environment, of the beam and generated fields on two space-separated vehicles—the main satellite APEX Interkosmos-25 and its subsatellite MAGION-3 (**Figure 1**). The independent subsatellite with small gas thrusters for orbit correction carried out simultaneous measurements on various mutual distances (from 10 m up to 1,000–2,000 km), at different areas of beam-disturbed environment, at different zones of the Earth's magnetosphere and ionosphere. The “boomerang” maneuvers of MAGION-3 were performed few times: it was pushed forward along the orbit and returned backward to main satellite, so MAGION-3 escaped from the main IK-25 up to few hundreds km and came back at distance about 0.4 km (**Figure 2**). An important part of the program were investigations of natural and human-made phenomena in a passive mode. The presence of two spacecrafts, equipped with practically equivalent complexes of the scientific equipment, enabled not only to carry out diverse researches of the ionosphere and the bottom magnetosphere, but also to distinguish spatial and temporal structures of the observable phenomena. A nearly polar orbit allowed to carry out research in auroral area. The complex program of “under satellite” ground observation and experiments included standard geophysical monitoring and special programs with ionospheric heating facilities.

Both spacecrafts, the main IK-25 and its subsatellite MAGION-3, were launched on an elliptical orbit with apogee of 3,080 km and perigee of 440 km, with the 82.5° inclination of the orbital plane on December 18, 1991 from the Russian space base Plesetsk. Orbital period was ~ 2 hours.

2.1. APEX Scientific Goals

The scientific goals of the project as described in Oraevsky et al. (1992) and Oraevsky and Triska (1993) are summarized below:

- Study of the dynamics and relaxation of modulated electron and plasma jets artificially injected into near-Earth plasma.
- Study of the low-frequency and MHD waves generation by heavy ions in the ionosphere.

- Determination of the radio emission generated by modulated beams of charged particles.
- Search for non-linear wave structures of the electromagnetic soliton type in a disturbed plasma environment.
- Study of the process of neutralization and dynamics of spacecraft charging.
- Study of the nature of the electrodynamic relationship between electromagnetic waves in the magnetosphere and ionosphere.
- Study of optical and radio emission in the auroral region.
- Passive mode observations—ionospheric plasma profiles under different conditions, auroral precipitation dynamics, mapping, and dynamics of polar cusps etc.

The APEX programme principal investigator (PI) was Viktor N. Oraevsky (1935–2006) from IZMIRAN, the MAGION subsatellite programme founder and PI was Pavel Triska (1931–2018) from GFI.

2.2. Scientific Payload

The main satellite IK-25 design was based on the AUOS-Z-AP-IK spacecraft bus aboard of the Tsyklon-3 rocket launcher, both produced by the KB Yuzhnoe, Dnepropetrovsk, Ukraine. The MAGION-3 subsatellite platform was designed and produced by the Geophysical Institute in Prague.

Both satellites were equipped with a complex scientific payload developed in a wide international cooperation. The particle, fields, wave, and supporting diagnostics are enumerated in **Tables 1, 2**.

2.2.1. IK-25 Short Description

The AUOS-Z bus consisted of a pressurized cylinder of about 1 m diameter and height ~ 2 m. Eight solar panels (~ 12 m²) were deployed 30° away (petal-like). The IK-25 satellite was gravitationally stabilized, its vertical position being kept by a spherical weight on about 20 m long boom (see **Figure 1**). The spacecraft utilized the 28 V power bus with the + pole commutation. Spacecraft charging characteristics (see section 3.1.1) suggested the solar panels were partially exposed to the ionospheric plasma. The current collection through the solar panel area in contact with the plasma depends on the panel design and technology—individual solar cell interconnects and edges exposure, coverglass conductivity/presence of a grounded transparent conductive coating, potential bias between the solar cells/cell interconnect layout etc., but such details of the AUOS-Z bus are unknown to us. For an example of the recent solar panel current collection modeling (see, e.g., Hess et al., 2016).

The UEM-2 electron gun consisted of two electron injectors G1 (oriented opposite to the satellite orbital velocity vector) and G2 (inclined toward the Earth). Utilizing a magnetic deflector, the G2 injector could turn the electron beam to the G1 beam orientation (see **Figure 3**) that corresponded to the injection pitch angles 50 – 80° for typical active parts of the IK-25 orbit. The electron acceleration voltage could reach 8–10 kV (unstabilized) with the peak current ~ 30 , 70, or 100 mA. The initial beam width was 4 mm at the output aperture. The electron gun could run in two basic modulation modes:

- Fixed frequencies: F40K (40 kHz), F15.6K, F976—all 2 μ s pulses, F3906, F30.5 (32 μ s pulses)
- Sweeping frequency: S250 (30.5 Hz–250 kHz, 2 μ s pulses, 1 DC + 11 frequency steps—1 s each) and S15.6 (30.5 Hz–15.6 kHz, 32 μ s pulses, 1 DC + 11 frequency steps—1 s each), basic cycle length 23 s

The UPM xenon plasma injector utilized a stationary plasma thruster (SPT) of Hall type (with close electron drift and extended acceleration zone) (Artsimovich et al., 1974) similar as used in the PORCUPINE experiment (Haerendel and Sagdeev, 1981). The xenon plasma was released to a 60° wide cone oriented $\sim 45^\circ$ from the zenith direction as illustrated in **Figure 3**. The ion energy reached 200–250 eV, ion temperature ~ 50 eV, ion current 2 A (10^{19} ions/s), $\sim 100\%$ ionization degree. As only xenon ions can leave the main discharge and accelerator section of the thruster, electrons from a supplementary xenon discharge (two external hollow cathodes LaB₆ installed, electron temperature ~ 2 eV) initiated the main discharge and simultaneously neutralized the xenon plasma beam injected to the space. The injector could work in these modes:

- Neutral xenon release, ~ 3 mg/s,
- DC xenon plasma emission–F0
- Modulated xenon plasma emission (almost 100% current modulation)–F1000 (1 kHz; 62, 125, 250, and 500 Hz modulation frequencies were other options)

The UPM cycle consisted of ~ 35 – 45 s preheating interval (with neutral xenon release lasting ~ 5 – 6 s, further denoted as Xe-A, and the hollow cathode emission, further as Xe-B) and one or several intervals of ~ 100 s plasma injection followed by ~ 50 s neutral xenon release only (Xe-A/Xe-B). Housekeeping and science data revealed additional ~ 0.12 Hz current modulation (up to 20%) due to unknown internal/external feedback. Only the neutral xenon gas was released during certain orbits as the ionization voltage was not applied due to technical reasons (Oraevskii et al., 1999).

Mean injection current profiles of the UEM-2 and UPM devices are displayed in **Figure 4**. The UEM-2 and UPM working cycles were not synchronized, one active experiment lasted usually ~ 3 – 5 min. Emission current and acceleration voltage (housekeeping) data were averaged over 0.32 or 0.08 s intervals. The UEM-2 and KM-10 floating probe telemetry data indicate a low intensity electron beam was also present during the one-second gaps between the modulated beam injections and in time between the S250 modulation sequences.

The KM-10 cold plasma monitor (also described in Afonin et al., 1994) consisted of the sensor and main electronics boxes. The sensor box was mounted on a ~ 0.8 m boom in front of the solar panels in the ram direction, out of the xenon plasma/electron beams. Its box with a conductive surface was electrically isolated from the spacecraft and kept near the ionospheric plasma potential actively by means of monitoring a floating potential (Φ_p) of one of its 7 probes (0.16, 0.64/0.2 s time resolution for the RTS/STO-AP telemetries, respectively, -10 to $+16$ V range). Ion density N_i , temperature T_{ix} , T_{iz} (planar retarding potential analyzers), ion drift in the YZ plane, and

TABLE 1 | IK-25 scientific payload (Dokukin, 1992).

Experiment	Function
UEM-2 electron accelerator	Injection of modulated electron beam (30 Hz–250 kHz), more than 200 injection sessions
UPM plasma gun	Injection of a Xe plasma beam, modulated current (DC, 32 or 1,000 Hz), more than 250 injection sessions
KM-10 cold plasma analyzer	Measurement of electron temperature (3 components), and ion drift (2 comp.), energy distribution, density, and temperature (2 comp.), floating potential, 50 ms resolution
NAM-5 mass spectrometer	Measurement of ion composition (1–60 a.m.) of plasma by radio frequency analyses, 3–4 ion species in 0.1 s, 8–10 ion species in 0.3 s
PEAS hot electron and ion analyzer	Measurement of a energy spectrum of electrons 30 eV–30 keV in 2 planes (12 angular sectors) during 1 s
DANI cold plasma complex	Measurements of a spectrum of energies and pitch-angular distribution of low energy electrons and ions (electrostatic analyzers) and parameters of plasma (Langmuir probe, Rogovsky belt)
NVK-ONCH low-frequency wave complex	Spectra of electrical and magnetic components of low-frequency electromagnetic radiation (0.1–20,000 Hz)
VCH-VK (PRS-3) high-frequency wave complex	Spectral analysis of high-frequency electromagnetic radiation, wave forms (0.1–10 MHz)
DEP-2 electric field experiment	Measurement of DC electric field, 3 components, 0.01–22 kHz
SGR-5 flux-gate magnetometer	Measurement of DC magnetic field, 3 components, $\pm 64,000$ nT, 0.64 (0.16) s resolution
ADO/service magnetometer	AUOS bus service magnetometer (attitude control), 0.64 (0.16) and 2 s resolution data available
MNCH-2 low-frequency search coil magnetometer	Measurement of low-frequency fluctuations of magnetic field, 3 components 0.1–10 Hz, 40 or 400 nT, sensitivity 1 nT
FS photometer	Scanning spectrometer, 3,194–6,563 Å
UF-3K 3-channel photometer	3-channel photo-spectrometer, 8 viewing sectors, lines: 3,194, 5,577, 6,300 Å

TABLE 2 | MAGION-3 scientific payload (Triska et al., 1990).

Experiment	Function
SGR-7/6 flux-gate magnetometer, magnetic variometer	Measurements of DC magnetic field, 3 components ($\pm 50,048$ nT) at frequencies 0–20 Hz and 1 component (± 159.4 nT) at 0.1–20 Hz with a sensitivity 50 pT
KEM-1 ELF/VLF wave experiment	Measurements of quasi-static electric field (3 components) 0.01–20 Hz, electric field (2 components) 0.1 Hz–120 kHz, magnetic field (1 component) 10–20,000 Hz, waveforms and filter banks / spectrum analyzer
KM-12 cold plasma analyzer	Measurement of temperature and density of cold plasma for 0.1 s
ZL-A-S Langmuir probe	Electron parameter fluctuations
PRS-2-S radio-spectrometer	Spectrum of electromagnetic waves at 0.1–10 MHz for 0.2 s/spec.; fluctuations of magnetic field ($f < 15$ kHz)
DANI-S low-energy particle analyser	Measurement of electrons and ions 0.05–1.0 keV for 0.1 s
DOK-A-S energetic particle spectrometer	20 keV — 1 MeV electron and ion fluxes, 8 energy levels, 0.1 s/spec.
MPS/PPS suprathermal particle analyser	Energetic spectrum of electrons 0.05–5(20) keV, 16 steps, 0.8 s/spec., 2 directions; energy and angular spectrum of electrons/ions 0.02–5(20) keV, 16 steps, 6 angular sectors;
FDS photometer	Measurement of optical ionospheric plasma emissions (6,300, 5,577 Å)

electron temperatures T_{ex} , T_{ey} , T_{ez} (planar probes with RF bias) were the output parameters of this instrument.

The PEAS hot plasma spectrometer (Nemecek et al., 1993, 1997) measured energy-angular distributions in two planes (2 double toroidal analyzers, 12 angular sectors each, electron/positive ion sweeps alternated). Its default operational mode provided one 16-energy level spectrum (50 eV–22 keV) of electrons and ions in 2.6 s (STO-AP), the sampling was not synchronized with the UEM-2/UPM data sampling. The AP-1 sensor block was mounted on a one-meter boom ahead of spacecraft (45° to XZ plane as depicted in **Figure 1**), the AP-2 sensor block was assembled ~ 0.5 m above the nadir spacecraft side (XY plane), the two block distance was ~ 3 m. Registered particle

pitch-angles (PA) distributions were computed using the ADO magnetometer. The block orientation corresponded to wide PA-range distributions by AP-1 (AP-2) in auroral (equatorial) regions, respectively. During flight, several sector channels got noisy and were excluded from further data processing.

The onboard wave measurements were performed using the PRS-3 (VCH-VK complex) plasma radio spectrometer, which represented a receiver with an input signal sensitivity of 0.5 μ V and a stepped tuning in the 0.1–10.0 MHz frequency range (Izhovkina et al., 2009). The frequency tuning step was 25, 50, 100 kHz, the bandwidth at a receiver input was 15 kHz, and the dynamic range of input signal level variations was 80 dB. An electric dipole antenna

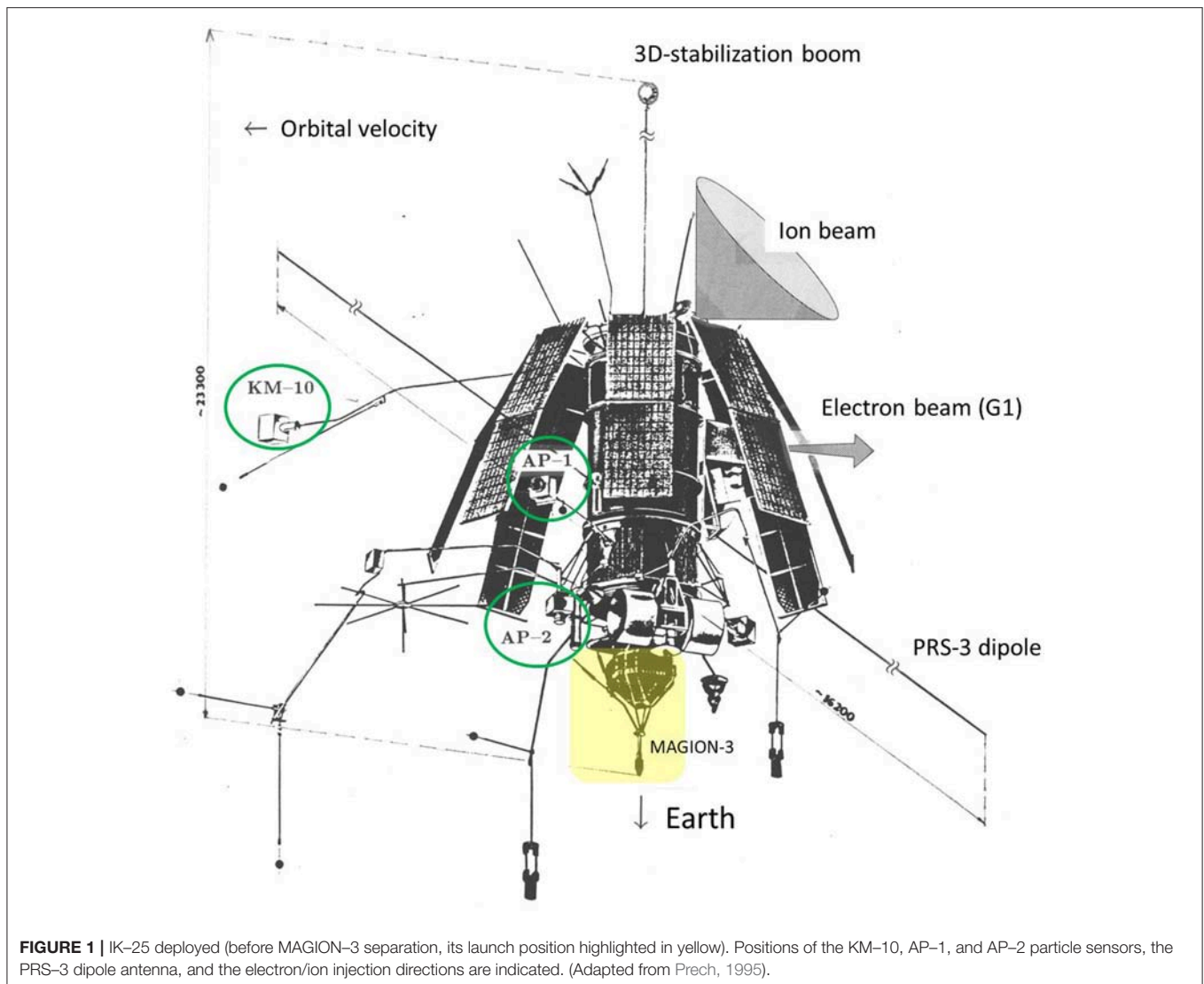


FIGURE 1 | IK-25 deployed (before MAGION-3 separation, its launch position highlighted in yellow). Positions of the KM-10, AP-1, and AP-2 particle sensors, the PRS-3 dipole antenna, and the electron/ion injection directions are indicated. (Adapted from Prech, 1995).

of the VCH-VK complex with a total length of 15 m, parallel to the Earth's surface, was used as a device sensor (Figure 1).

The low-frequency wave instrument NVK-ONCH registered magnetic field spectrum in the range 8–969 Hz and amplitudes at fixed frequencies 9.6 and 15.0 kHz (Oraevsky et al., 2001). Measurements of the sensitive flux-gate magnetometer SGR-5 and the spacecraft service magnetometer were used to assess a level of magnetic field fluctuations at lowest frequencies. Quasi-steady electric field 0.1–10 Hz components and VLF electric field spectra were measured by a system of double electric probes connected to inputs of the DEP-2 instrument and the NVK-ONCH wave complex (Baranets et al., 2007).

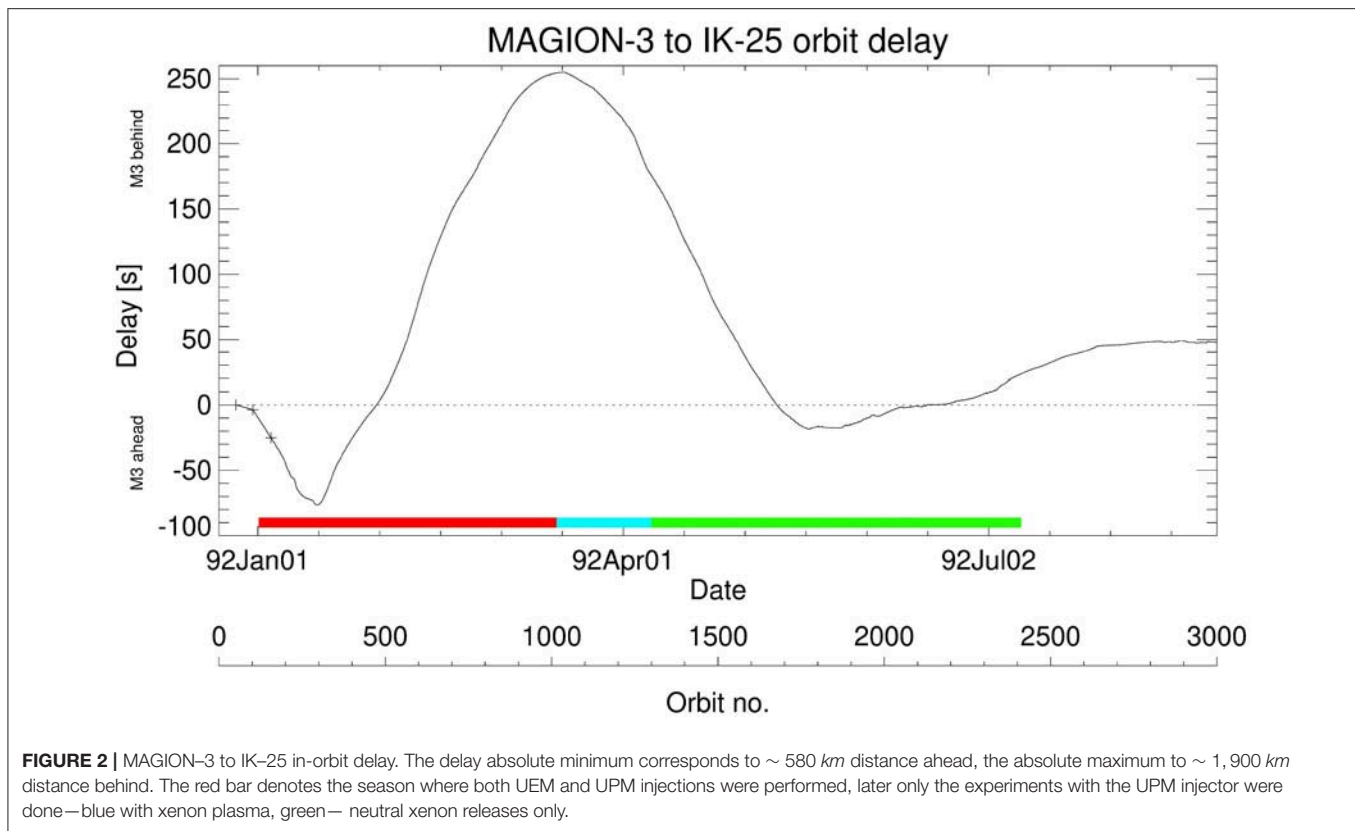
2.2.2. MAGION-3 Short Description

The size of the object was 0.85×0.60 m diameter (2 m with deployed booms), its weight was 52 kg. The orbital

speed could be corrected using pressurized neutral gas (nozzles parallel/antiparallel with the satellite axis). The satellite axis was oriented approximately along the local geomagnetic field (built-in permanent magnet 30 Am^2 with week dumping) (Triska et al., 1990).

The suprathermal particle spectrometer MPS/PPS was a simplified version of the PEAS instrument (Nemecek et al., 1994). It measured energy distribution of electron and ions in one half-plane containing the satellite main axis, divided to 6 angular sectors (SEA-A and SEA-B toroidal-cut analyzers, 25° width). The system was complemented by two narrow field-of-view electron energy analyzers MP-A and MP-B (parallel/antiparallel to the satellite axis, respectively) to obtain a full-range electron pitch-angle distribution (8 channels with 30° spacing).

The energetic particle sensor DOK-A-S registered energetic electron and proton fluxes parallel and perpendicular to the satellite axis using two pairs of silicon detectors (Prech et al., 2002; Baranets et al., 2007).

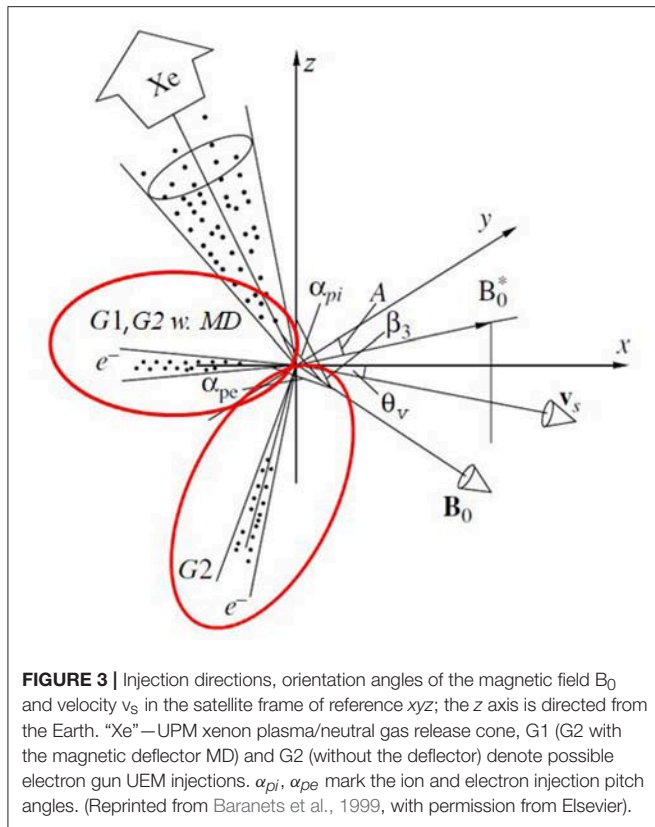


The wave experiment provided spectra of two electric and one magnetic field components in ELF/VLF ranges (Triska et al., 1990; Baranets et al., 2007). Waveforms of selected electric/magnetic components up to 60 kHz ($2-3 \times 8$ Hz–20 kHz or narrow-band) could be directly transmitted to ground. The PRS-2-S radio-spectrometer provided spectra with Δf 15 or 50 kHz resolution in the HF frequency range 0.1–10 MHz (one electric field component, 3 m length dipole) (Rothkaehl and Klos, 1996).

2.3. Active Experiment Methodology

The MAGION-3 was separated on December 28, 1991 from the main satellite and drifted away with a speed ~ 5 km/day. Several tens of orbit corrections were made to keep the mother-daughter relative distance within range -580 to $+1,900$ km. The in-orbit MAGION-3 to IK-25 delay evolution is depicted in Figure 2. Color bars in this figure mark different campaigns of active experiments, planning of which was affected by the electron and ion injectors working status. About 200 active experiments with electron and plasma injections in various configurations were performed till July 1992 within the satellite relative distance 70–500 km. Planned experiments in near (> 10 m to 0.4 km) and mid-range ($\sim 1-10$ km) zones could not be performed for technical reasons. Two-point passive measurements at distances 100–2,000 km continued till the MAGION-3 end-of-life in August 1992. Ionospheric investigations by IK-25 lasted till July 1993.

The IK-25 satellite was controlled from Soviet telemetry stations and selected scientific data were delivered to ground via its primary telemetry channel (RTS), mostly from its on-board telemetry memory with downloads once per day but real-time telemetry sessions were also started frequently in the regions of satellite visibility to allow better data time resolution. Except the active experiment devices and diagnostic equipment IK-25 also included a complementary telemetry system STO-AP developed by teams from Hungary, former USSR, Poland, and Czech Republic. The system STO-AP provided an interface for the scientific instruments, management of operation modes, preliminary processing of the scientific information, independent control of onboard experiments, formation of the TM frames and transfer of the telemetry information to the ground-based stations at Czech Republic (Panska Ves), Russia (Troitsk, Apatity, Tarusa), and Germany. STO-AP enabled higher volume/better time resolution of science data mostly during real-time telemetry sessions lasting about 10 min (memory replay sessions with limited data volume were also sporadically performed). Drawback of this approach was the IK-25 visibility regions from the RTS and STO-AP reception stations did not fully overlap (resulting many active experiments controlled via RTS were not fully covered by scientific data routed via STO-AP). Beside it, the STO-AP telemetry channel (137/400 MHz bands) was often affected during strong xenon plasma injections (short drop-outs and wide data gaps were present during such sessions). The limited total telemetry capacity



necessitated trade-offs in the scientific instruments operational modes and their resolution, during some orbits not a complete parameter set could be studied.

The MAGION-3 daughter satellite was controlled and scientific data transmitted via the STO telemetry to the Panska Ves observatory (Czech Republic) mostly during real-time sessions as its onboard data memory volume was limited (4 MB). For this reason the satellite pair visibility also affected the simultaneous data coverage of active experiments as performed from the main IK-25 satellite, especially during periods of larger spacecraft separation.

The Panska Ves visibility requirement made the altitude, invariant latitude (INL)/L-shell, magnetic local time (MLT) parameters of the APEX active experiments bound. As the active experiments were run above the north hemisphere, downward electron injections during day were toward the north magnetic pole (near mirror point in dense ionosphere below) and upward electron injections (during night) were directed toward the south magnetic pole (distant mirror point behind the rarefied top ionosphere/magnetosphere). This brings difficulties to separate some physical parameter dependencies. The electron and plasma injectors onboard IK-25 had many working/modulation modes and the elliptical orbit of the satellite pair allowed performing the active experiments for different ionospheric parameters. Ideally, a certain injection configuration should be repeated several times under different ambient conditions, but in reality only one usable dataset was captured for many configurations.

Despite these technical difficulties a rich database of observational data from active experiments of the APEX project was collected and many interesting scientific results were published.

3. RESULT OVERVIEW

3.1. Spacecraft Charging and Neutralization During Active Emissions

Beside the complex IK-25 design geometry the evaluation of the spacecraft (s/c) charging properties was complicated also by other factors. The IK-25 body was painted and its surface was assumed equipotential but the conductivity was not characterized in publicly available data. The scientific payload was painted or covered by multi-layer insulation (MLI) and connected to the spacecraft ground and the surface again assumed equipotential (not guaranteed). Solar panel surfaces probably were not equipotential—the positive potential end was probably exposed to the ionospheric plasma (dayside effects on the satellite potential). Also, after the end of all injections the spacecraft potential return to its “quiet” level immediately during night, but with a time constant of few minutes during day. A full charge balance analysis has also to consider the stabilization boom/weight at ~ 20 m distance. The electrical current across the boom was not monitored, unfortunately.

The KM-10 sensor box (placed ahead of the spacecraft in the ram direction) was kept floating near the ionospheric plasma potential ($\delta\Phi \sim -0.2$ to -0.7 V for $T_e \sim 0.1 - 0.3$ eV, not considering photocurrent/suprathermal electrons). The KM-10 floating probe potential Φ_p (normally positive against s/c) was widely used as a spacecraft potential proxy ($\Phi_s \approx -\Phi_p + \delta\Phi$), but its design range ± 90 V is doubtful (-10 to $+16$ V limitation was observed in practice). The KM-10 data time resolution (0.2 s) and UEM/UPM injection data (0.32/0.08 s resolution) unfortunately do not allow to study transient charging phenomena during the injections on/off or detail profiles for individual sub-millisecond electron injection pulses. As the Debye length near the apogee (~ 10 V, 10^4 cm $^{-3}$) is approximately 1 m, the KM-10 might not be always out of the spacecraft Debye sheath. Also, this spacecraft potential proxy is considered unreliable when the KM-10 ion density N_i drops to its lower limit ($\sim 10^8$ m $^{-3}$). The parameter reliability was discussed in more details in Prech (1995).

During the active experiments with electron injections, the secondary electron emission from the spacecraft surface due to hot electron collection (return currents) probably also affected the spacecraft charge balance. Up today no detail and realistic IK-25 charging model is publicly available. A simplified model is discussed by Zilavy et al. (2003).

3.1.1. IK-25 Potential Outside of Active Experiments

Prech et al. (1999) investigated the “quiet time” behavior of the IK-25 spacecraft properties. Figure 5A depicts a potential difference between the KM-10 floating probe and the satellite (Φ_p) as a function of the altitude. The data were collected during intervals preceding the electron/plasma injections and the satellite potential can change for several volts depending on the part of the orbit (with changing the altitude, ambient plasma

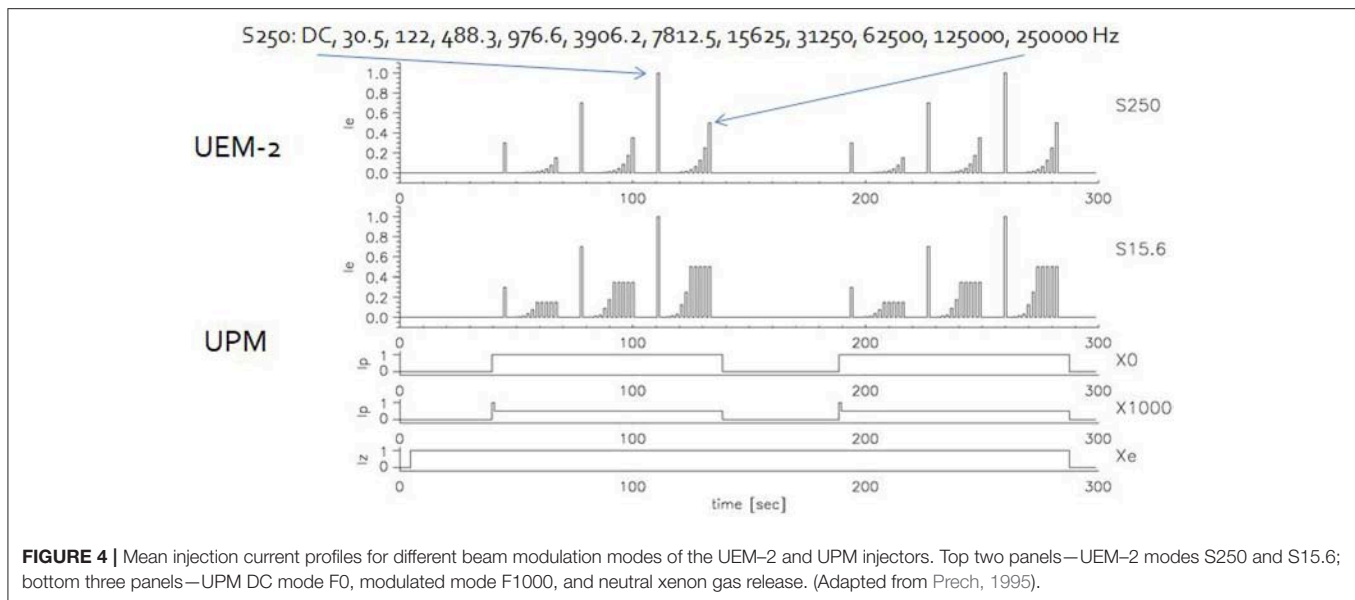


FIGURE 4 | Mean injection current profiles for different beam modulation modes of the UEM-2 and UPM injectors. Top two panels—UEM-2 modes S250 and S15.6; bottom three panels—UPM DC mode F0, modulated mode F1000, and neutral xenon gas release. (Adapted from Prech, 1995).

density, and electron temperature), the satellite always charges negatively. The dayside/nightside orbits split the observations to two branches. The night branch (the satellite was in the Earth's shadow) is a rising function of the altitude which was explained with the electron temperature also rising with the altitude. **Figure 5B** shows the same data as a function of the ion density N_i measured by the KM-10 device. The higher potentials (more negative s/c charging Φ_s) on the dayside branch are assumed due to the solar panels connected through a small resistance to the ambient plasma at low altitudes. This resistance is a function of density and the effect gradually disappears at altitudes above 2,500 km or densities below 10^9 m^{-3} .

3.1.2. IK-25 Potential During the Electron Beam Injections With Xenon Plasma Neutralization

The IK-25 spacecraft charging during the electron beam injections with the xenon plasma neutralization can be summarized to following conclusions (Prech, 1995, 2002; Prech et al., 1995, 1999; Nemecek et al., 1997):

- For ambient plasma densities $N_i > 10^{11} \text{ m}^{-3}$ the electron beam charge was well neutralized, $\Phi_s < 15 \text{ V}$ for S250 and F40K modes peak currents, KM-10 density N_i measured ahead of the spacecraft was not further increased during the xenon injections.
- For ambient plasma densities $N_i < 10^{11} \text{ m}^{-3}$ and the xenon plasma release during day kept the KM-10 ram density $N_i \sim 8 \times 10^{10} \text{ m}^{-3}$ and the electron beam charge was well-neutralized, the neutralization appeared even better during night, the KM-10 ram density $N_i \sim 1.5 \times 10^{11} \text{ m}^{-3}$ and the spacecraft potential relaxed to its “quiet” level (for example see **Figure 6B**).
- For lower ambient plasma densities during neutral xenon gas / hollow cathode releases (Xe-A/Xe-B) the spacecraft is charged to more than +20 V against its “quite” level.

The KM-10 floating potential reached a negative saturation or after shortly negative excursion it traveled to positive saturation and as such it could not be used as a reliable IK-25 charging monitor (**Figure 6B**). As accelerated cold ionospheric electrons were not registered by the PEAS spectrometer, the spacecraft potential was within limits $\sim 20 \text{ V} < \Phi_s < \sim 50 \text{ V}$.

- During the modulated electron beam injection (S250), when Φ_p was not in saturation, $\sim 1 \text{ s}$ positive pulses in the spacecraft potential were observed (rising/falling edges $< 0.2 \text{ s}$)—see e.g., **Figure 6A**. Their amplitude increased with the mean electron pulse current level until it reached the unmodulated (F40K) levels. The amplitude was larger during Xe-B mode than with the main xenon plasma injection showing worse electron beam neutralization in this mode.

3.1.3. IK-25 Potential During the Xenon Plasma Injections

The IK-25 spacecraft potential behavior during the UPM neutral gas releases and xenon plasma injections (UEM not active) was studied by Prech et al. (1999), further examples and discussion are contained in Prech (1995, 2002). It was found that the release of the neutral xenon alone did not affect the spacecraft potential under conditions of the APEX experiment. The change of the spacecraft potential $d\Phi_p$ induced by the release of the xenon plasma was up to $\sim \pm 10 \text{ V}$ (e.g., **Figures 7A,B**). Such a change was considered reasonable, as the temperature of the plasma in the injector output was about 50 eV, sharply decreasing with the distance from the spacecraft due to adiabatic expansion of the xenon plasma cloud. A dense conducting cloud (plasma contactor) created by the released plasma behind the spacecraft was expected to effectively collect electrons from the ambient plasma. Being partly in the spacecraft wake the cloud contribution to the ion current is found smaller. The dimension of this cloud was expected comparable with the dimensions of the “diamagnetic

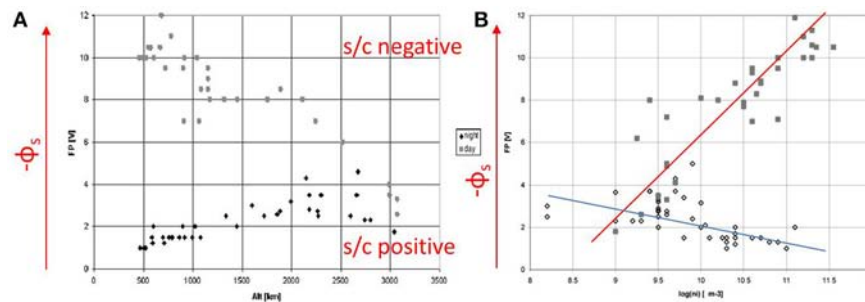


FIGURE 5 | IK-25 spacecraft potential outside of active experiment intervals. The plots show the potential difference between the KM-10 floating probe and the spacecraft body as a function of the altitude **(A)** and the ambient plasma density **(B)**. The spacecraft is charged negatively for positive KM-10 floating potentials. (Reprinted from Prech et al., 1999, with permission from Elsevier).

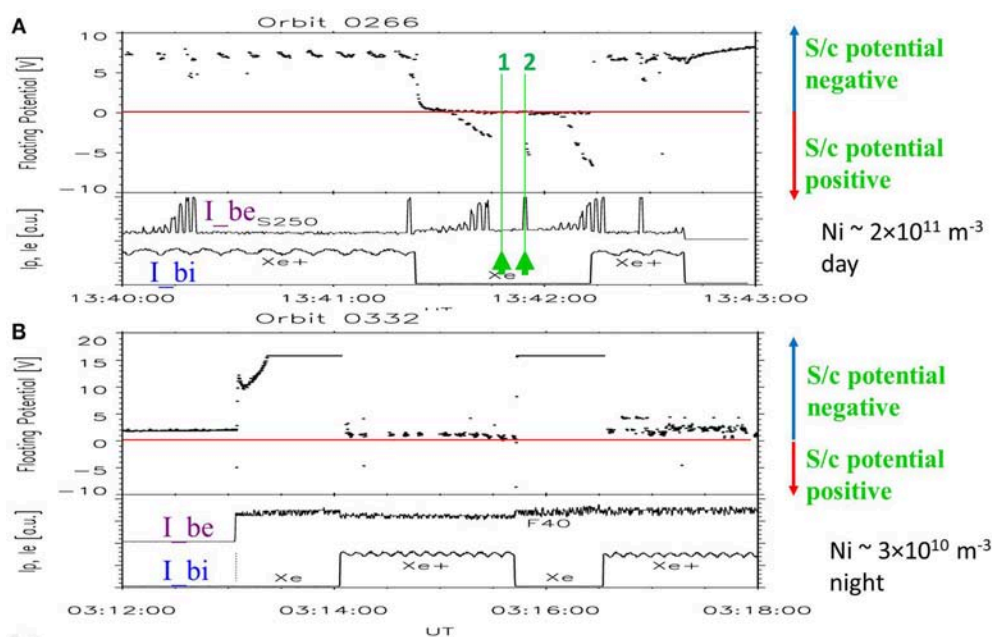


FIGURE 6 | Examples of the IK-25 spacecraft potential behavior during electron beam injections. The plots show the KM-10 floating probe potential, mean UEM electron beam and UPM plasma injector currents: **(A)** Orbit 0266—dayside modulated electron beam (S250 mode, G1 injector) with xenon gas/plasma neutralization (F0 mode), green vertical lines assigned 1 and 2 mark instants with the UEM beam base current and UEM beam maximum current (DC pulse); **(B)** Orbit 0332—nightside electron beam injection (F40K mode, G2 injector with magnetic deflector) with xenon gas/plasma neutralization (F0 mode). (Reprinted from Prech et al., 1995, with permission from Elsevier).

phase” of the plasma cloud expansion proposed by Hausler et al. (1986).

The electron current collected by the cloud should rise linearly with ambient plasma density, but this rising current causes an increase in the negative spacecraft potential. The result was the observed linear relationship between $\log N_i$ and dFP (Figure 7C). The collected electron current as the difference between those electrons actually collected and those leaving the plasma cloud was found either positive or negative. If the number of electrons collected from the ambient plasma is higher than the number leaving the cloud then the spacecraft potential change vs. “quiet” level dFP is

negative, otherwise it is positive. Altitude, ambient plasma temperature, and day/night dependencies were found small for this effect.

The DC xenon plasma emission was studied in a simple model of a planar floating probe, spherical satellite, and plasma cloud by Zilavy et al. (2003). The authors stress the importance of suprathermal electron tail and ion mass composition (affecting the H^+ , O^+ ion ram energy) of ionospheric plasma for charge balance. The model can explain the spacecraft potential change of both polarities suggesting the ionospheric O^+ ions are less effectively collected by the plasma cloud due to their higher ram kinetic energy.

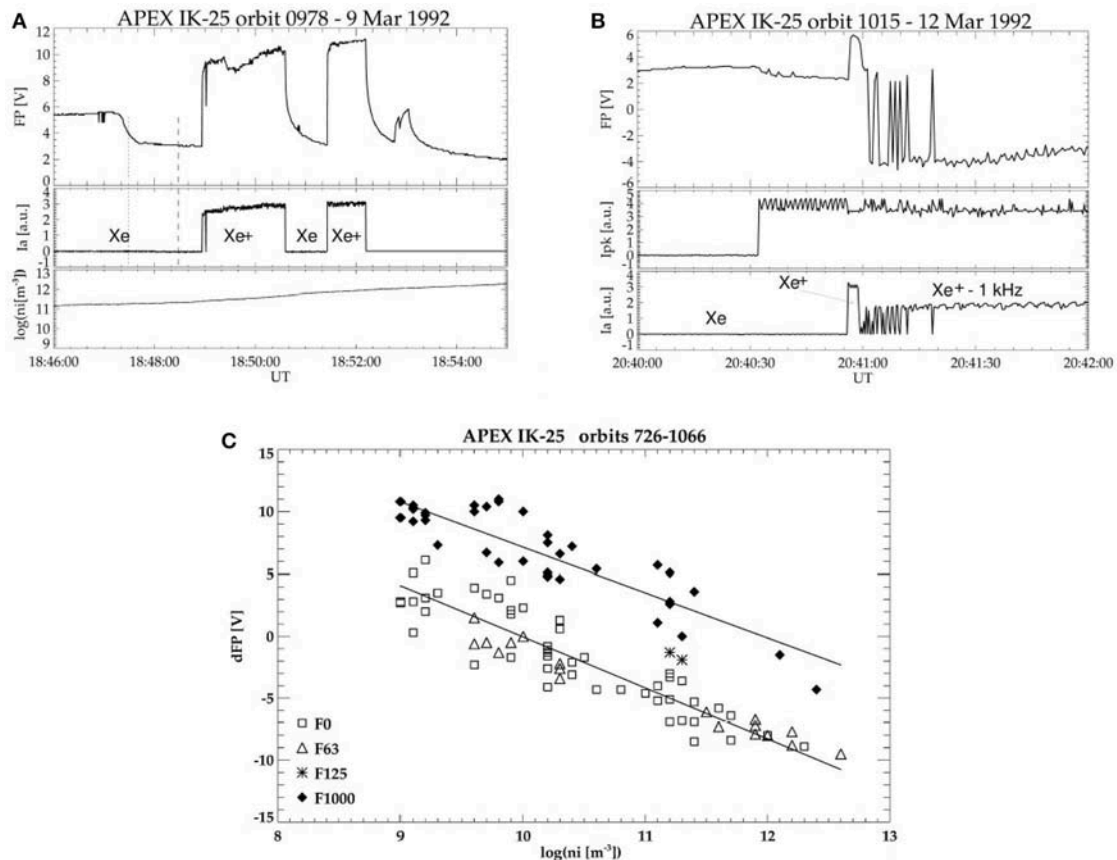


FIGURE 7 | Examples of the IK-25 spacecraft potential behavior during the xenon plasma injections. **(A)** (Night DC injection F0 mode at ~ 440 km altitude) shows the KM-10 floating probe potential, the mean UPM ion current, and the ambient plasma density (KM-10). **(B)** (Dayside DC/1 kHz injection F1000 mode at $\sim 2,900$ km altitude) depicts the KM-10 floating probe potential, the UPM hollow cathode current, and the mean UPM ion current (in both panels Xe stands for the neutral xenon release intervals, Xe+ for the xenon plasma injections). **(C)** Change of the spacecraft potential during the xenon plasma releases as a function of the ambient plasma density for different UPM modulation modes. (Reprinted from Prech et al., 1999, with permission from Elsevier).

3.2. Neutral Xenon Release Experiments

During the xenon plasma injection experiments of the APEX programme, about 20 experiments were performed when the UPM device worked uniquely in the mode of neutral xenon gas release. Additional experiments included simultaneous injections of electron beam and neutral xenon. Under conditions of the experiment the kinetic energy of the xenon atoms motion relative to the ionospheric plasma ($\sim 25\text{--}40$ eV) exceeded the xenon first ionization potential (~ 12.1 eV), so the necessary condition of the anomalous ionization was satisfied. The neutral xenon releases occurred over the full altitude range 450–3,000 km and variety of experimental conditions (geomagnetic field strength, ambient plasma density, injection angle vs. magnetic field direction, illumination etc.) allowing comparison with other experiments.

Oraevskii et al. (1999) and Choueiri et al. (2001) analyzed the influence of the neutral xenon releases to spectra of high-frequency plasma turbulence, accelerated electron spectra, and spectra of electric field fluctuations near lower-hybrid resonance. They also studied dependence on angle between the magnetic

field and the gas injection direction. The experimental data show the ambient plasma response to the neutral gas injection including electron temperature and anisotropy increase and amplification of wave activity practically over full range of registered frequencies. Theoretical considerations of these papers are consistent with the observed data. Although only small flows of neutral xenon gas (~ 3 mg/s) were released during several minute lasting intervals, due to its collisional interaction with the background ionospheric plasma a sufficient number of “seed” energetic ions was created so as the CIV-related electron heating could be observed. Newly created xenon ions by charge exchange collisions, electron impact ionization, photoionization, and scattering of ionospheric plasma ions were considered in their analyses. The calculated ionization yield smaller than 1% corresponds to not observable changes of plasma density. Ion-ion lower-hybrid wave instability was expected to reach distances 3–100 m. In the spacecraft reference system these waves have a perpendicular to the magnetic field phase velocity component in the range of reflected ion speeds 0–7 km/s, but their group velocity is directed almost parallel to the

magnetic field. For this reason the waves can reach the spacecraft and be detected. Unfortunately, the proper frequency band diagnostic was not operating, but simultaneous narrow- and wide-band HF emissions in the range 3–10 MHz during dayside xenon gas emissions with pitch angles 85–115° were observed. Intensification of HF wave activity was not observed during nightside emissions (pitch angles 57–71°). The wave frequency bands were flat and did not evolve along the orbit (magnetic field strength). The authors think that the observed waves could be a symptom of turbulent fluctuations connected with the instability, but without a theoretical explanation. According to the authors, it is possible that the presence of the solar flux and not the emission pitch angle is the controlling parameter for the observed effects. In such a case, the role of the solar flux in affecting HF wideband activity through intermediary effects such as plasma enhancement due to photoionization may be worthy of investigation.

3.3. Beam-Plasma Interaction

3.3.1. Hot Electron Pitch-Angle Distributions During Electron Beam Injections

Plasma environment in the vicinity of the IK-25 spacecraft and return currents during the UEM electron beam injections were studied in Prech (1995), Nemecek et al. (1997), and Prech et al. (1998) using the PEAS electron and ion spectra with following conclusions:

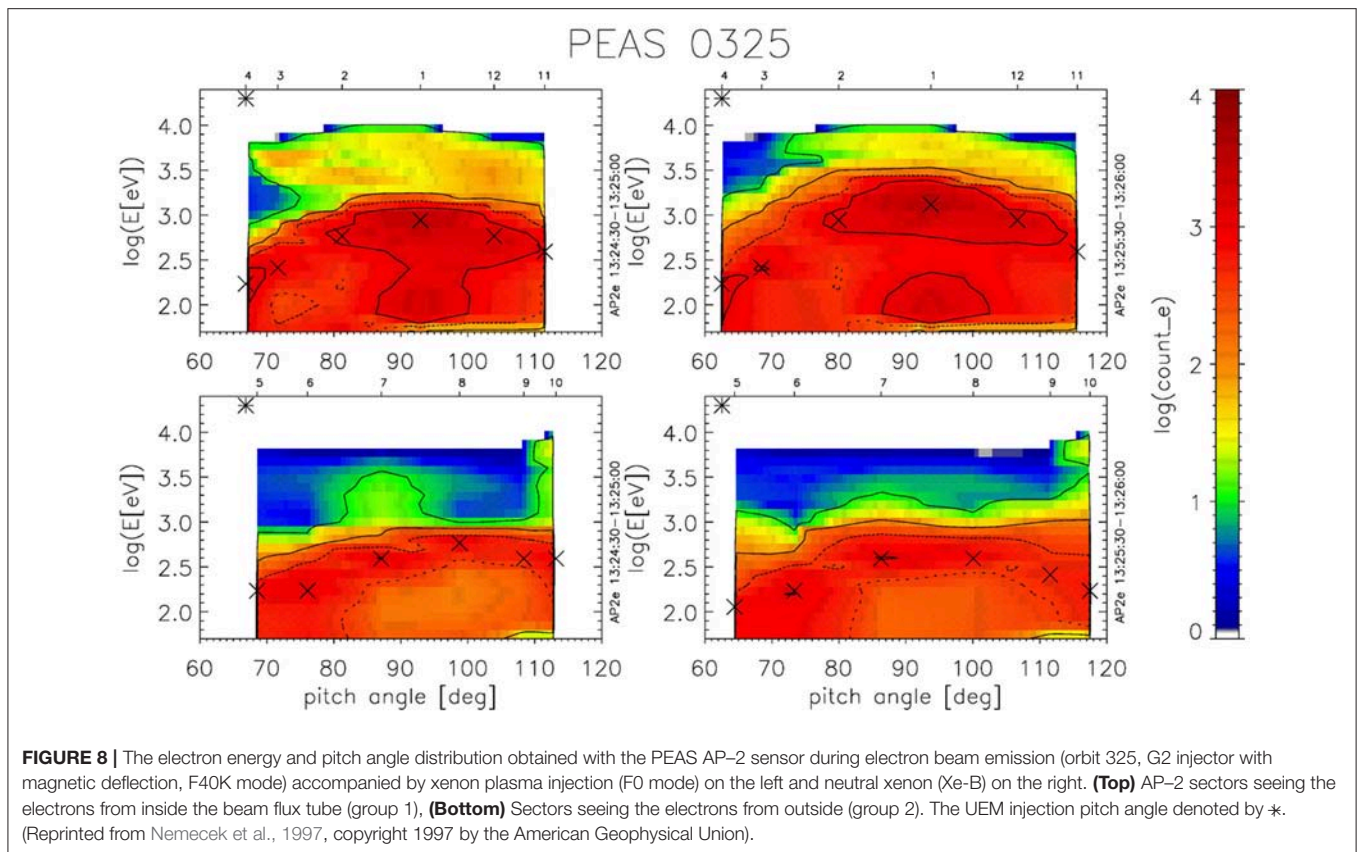
- Hot electrons as a part of the return current were registered by both the AP-1 and AP-2 sensors with energies peaking between $\sim 10^2$ and 10^3 eV, electrons were preferentially accelerated perpendicular to the magnetic field. The acceleration mechanism was weaker during simultaneous xenon plasma release.
- The hot electron flux and energy increased with the mean primary beam current, but the PEAS time resolution did not allow to observe any resonance effects depending on the electron beam modulation frequency.
- The energy/pitch-angle distributions allowed to distinguish several electron populations attributed to accelerated ambient plasma electrons (at lower energies) and scattered beam electrons (few keV energies). No electrons with energies accelerated above the primary beam energy were found in the PEAS spectra.
- Hot suprathermal electrons constituting a part of the return currents arrived with comparable fluxes from directions parallel and antiparallel to the magnetic field, i.e., not only from the area behind the spacecraft disturbed by the injected beam but also from the ahead region that could be affected only by the return currents and secondary products of the beam-plasma interaction.
- The energy and flux of these electrons decreased for pitch angles closer to parallel/antiparallel directions indicating the space localization of the acceleration process.
- The electron distributions were not homogenous in space—the AP-1 sensor registered distributions narrower in energy and lower fluxes than the AP-2 sensor mounted close to the spacecraft body.
- The AP-2 electron spectra showed differences in pitch-angle spectra for sectors looking inside (group 1) and outside (group 2) of a cylinder shell containing the primary beam electrons (an example is presented in **Figure 8**). The asymmetry of the electron fluxes could be explained with $E \times B$ drift but the corresponding local electric fields would reach ≈ 100 V/m.
- The observations supported the picture where the primary beam created a hollow cylinder with its shell full of hot electrons, the cylinder extended both parallel/antiparallel to the magnetic field. The AP-2 sensor was located right in the hot shell. The physics of such beam expansion is, of course, complex and as such studied theoretically including computer simulations and experimentally both in laboratories and in space. One of the relatively recent computer simulations related to the APEX experiment was made by Lizunov et al. (2002). The beam current profile and generally the beam-plasma interaction in context of the APEX experiments was recently investigated e.g., by Baranets et al. (2012, 2017).

No significant ion fluxes were registered by the PEAS sensors during the UEM and UPM active injections (in different mode combinations), the counts usually remained near the noise level (Nemecek et al., 1997). Only sporadically an ion group with energy corresponding to the UPM acceleration voltage was registered in a narrow spatial angle—usually just one sector. Although the PEAS sensors had no mass resolution capability it was assumed these were ions from the plasma cloud edge that after one gyro-revolution returned to the PEAS input aperture.

More recently, Budko et al. (2003) also analyzed the PEAS electron and ion distribution behavior during the electron beam injections related to the HF wave emissions (see section 3.3.3). They emphasized detection of very short-term bursts of accelerated ions with energies up to several hundreds of electronvolts immediately after switching on the electron gun. The effect again corresponds to generation of strong electric fields in the near zone.

3.3.2. Plasma Environment in the IK-25 Spacecraft Vicinity, Diamagnetic Cavity

The UPM plasma cloud density evolution in the APEX active experiments were estimated by Volokitin et al. (2000). The kinetic energy 200 eV of accelerated xenon ions corresponds to ~ 17 km/s speed. For the typical UPM emission parameters 2.5×10^{19} ions/s (~ 3 mg/s) assuming density decreased as $\sim 1/R^2$ the xenon plasma expanded to the ambient plasma density levels $\sim 10^{10} \text{ m}^{-3}$ at distances $R \simeq 300\text{--}400$ m. From similar estimations the xenon plasma dynamic pressure exceeded the geomagnetic field pressure for $R[\text{m}] < 2.5 \times 10^5/B_0[\text{nT}]$. For the APEX orbit ($B \sim 3\text{--}5 \times 10^4$ nT) the plasma cloud could expel the geomagnetic field and create a diamagnetic cavity to distances 5–10 m from the spacecraft (the xenon ion gyroradius 1–2 km) due to electrical polarization of plasma cloud. While the injected energetic xenon ions move more-or-less freely across the magnetic field, the electrons are magnetized and follow the magnetic field lines. The electrons partly drift due to the polarization electric field across the magnetic field and at the same time they drag the magnetic field lines in their effort to



follow the cloud ions. However, this process could be noticeably depressed if the ambient plasma electron density is sufficient to support the parallel electric current neutralizing the polarization electric field. Volokitin et al. (2000) studied the diamagnetic cavity formation during the IK-25 xenon plasma injection using two magnetometers (see **Figure 9**). The SGR-5 magnetometer mounted on a boom at least 7–12 m from the xenon plasma beam registered magnetic field fluctuations with amplitude ~ 10 nT while the service magnetometer mounted very near the injected cloud observed magnetic depressions 1,500–2,000 nT. The experimental data documented the active extrusion of the geomagnetic field in the initial phase of the plasma cloud expansion. The authors compared the observational results with a theoretical model based on ideas from analysis of previous experiments PORCUPINE (Hausler et al., 1986; Oraevsky et al., 2003) and AMPTE (Mishin et al., 1988).

3.3.3. Beam-Plasma Interaction, ELF/VLF, and HF Waves Generation

Near-zone VLF wave observations (NVK-ONCH instrument) during the APEX xenon plasma injection have been studied by Mikhailov et al. (1998, 2000). While the unmodulated xenon plasma emission was accompanied by a broadband VLF noise, during the 1 kHz and 125 Hz beam current modulation (F1000 and F125 modulation modes), the basic frequency and its harmonics were registered in the magnetic field data. Simultaneous disturbances in electron temperature and plasma

density from the KM-10 device were detected. While the onset of fast magnetic field variations coincided with the plasma injection switch-on, they lasted much longer after the injection off, with time-varying spectral amplitude profiles across the ELF/VLF range. The authors attributed the observed delays to a joint movement of the satellite and a packet of waves created inside the plasma cloud (assuming different wave modes were excited during the injections and some of them were able to travel in the direction of the satellite velocity according to the authors). Similar results were obtained also in previous higher flow, ram directed, xenon neutral gas releases from the ACTIVE Interkosmos-24 satellite (Klos et al., 1998).

Kiraga et al. (1995, 1998) analyzed HF emissions during modulated electron beam injections (F40K mode, $2\mu\text{s}$ pulses, neutral xenon simultaneous release) in the middle latitude dayside ionosphere (altitude $\sim 700 - 1,100$ km) and near the nightside perigee (altitude $\sim 415 - 470$ km, middle and high latitudes). In the regime of strong beam-plasma interaction, the HF dipole antenna was effectively screened from the HF broadcasting stations radiation as probably a large volume of excited plasma around the emitting satellite scattered these waves. Complex structures of peaks near harmonics of local plasma, cyclotron, and upper-hybrid frequencies were present in the HF wave spectra of PRS-3 during the electron injections, varying with ionospheric parameters and evolving injection configuration along the orbit. Doublet of local plasma and upper hybrid frequency peaks was the most prominent

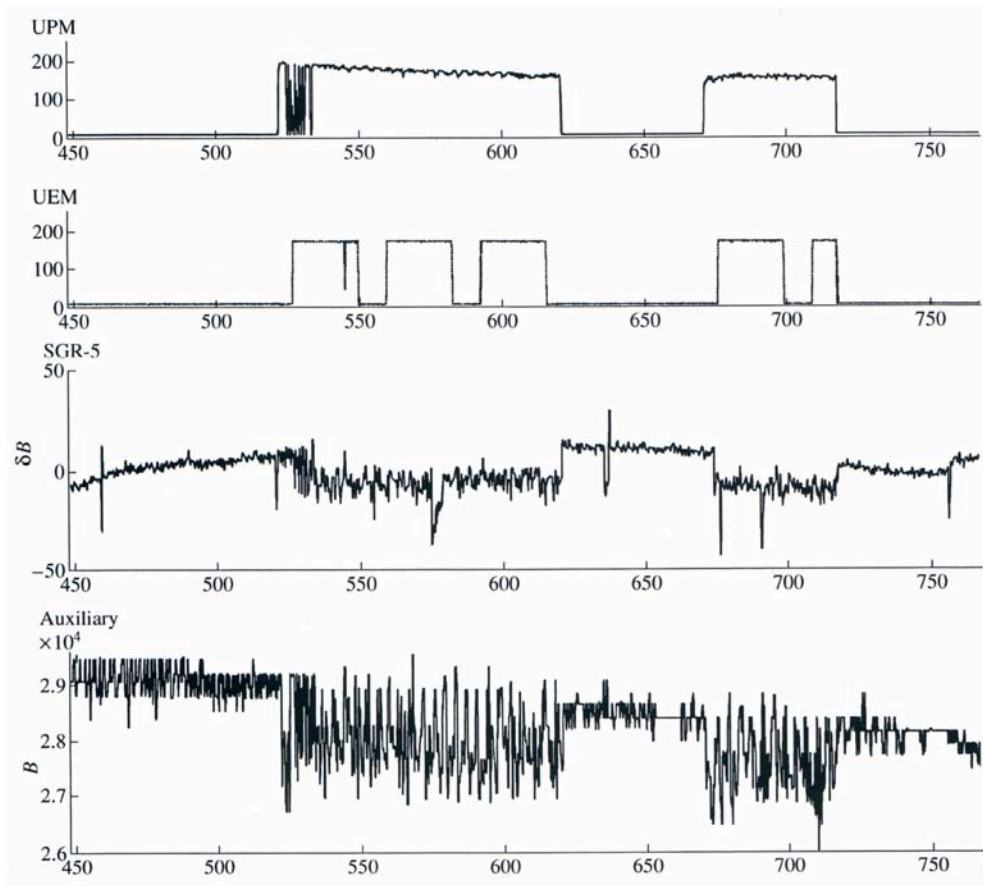


FIGURE 9 | Illustration of the diamagnetic cavity formation as seen in the time-dependent variations of the magnetic field. Horizontal axis—the time [s] lapsed from 04:37:02 UT, obtained during the xenon plasma injection on March 13, 1992 (orbit 1019). From top to bottom the UPM and UEM acceleration voltage [a.u.], deviation of the magnetic field strength [nT] from a mean level (SGR-5), and the service magnetometer total magnetic field [nT]. (Reprinted from Volokitin et al., 2000).

emission during the former interval and harmonics of this structure were also excited. During the latter injection the recorded spectra were more diverse, main emissions consisting of even harmonics of the cyclotron frequency f_c , dublets below $2f_c$, or selectively excited emissions near $3f_c$ and $4f_c$. The authors interpreted the individual spectral peaks as emissions characteristic for synchrotron maser radiation created in a system of cold magnetoplasma and cold, weakly relativistic, very diluted electron beam.

Considering the electron beam under the F40K modulation as “strong,” Kiraga (2003) brought a picture of HF emissions obtained during “weak” electron beam emissions in the S250 and S15.6 modulation modes. Evolution of spectra during the beam cycle, presence of modulation frequency harmonics and structures around the upper hybrid frequency are discussed as well as effects of antenna charging and electron current coupling between plasma and antenna. The electron beam injection as a tool for continuous monitoring of ambient plasma density by reception of emissions generated in ambient (f_p, f_{uh}) frequency band is also discussed in the paper.

Budko et al. (2003) investigated HF waves excited by the pulsed electron injections and xenon plasma releases in the

near zone. In lower altitude, subauroral dayside ionosphere they recorded emissions at the local plasma frequency f_p , $2f_p$, and $4f_c$ during a weak electron injection alone. Later on during the simultaneous electron beam (S250) and xenon plasma (F0) injections, whistler mode wave excited at the electron beam modulation frequency ~ 250 kHz and resonances at the first three cyclotron frequency harmonics were present in the PRS-3 frequency range (the local plasma frequency during the xenon plasma release was estimated by the authors to ~ 16 MHz which made detection of plasma resonance of an electron beam and background plasma impossible). During strong electron injection modulated by the 250 kHz frequency, with only neutral xenon release, the dublet of local plasma and upper hybrid frequency peaks (reported by Kiraga et al., 1995 for 40 kHz modulation) was clearly present together with whistler mode waves possibly up to the 6th to 8th harmonic of the electron beam modulation frequency. During nightside subauroral and middle latitude electron beam injection (F40K), resonant peaks near even electron cyclotron frequency harmonics ($2f_c, 4f_c, 6f_c, 8f_c$) or combinations with the plasma frequency ($f_p, 2f_c + f_p, 4f_c + f_p, 6f_c + f_p$) were recorded in the former case, plasma resonance harmonics only in the latter case.

Recently, very detail analyses of the beam-plasma instability development, excitation of VLF/LF and HF waves, coupling of waves, particle heating and acceleration related to different APEX experiment configurations (injections through/opposite the xenon plasma cloud, quasi-perpendicular electron beam injections, injections of differently modulated electron beam - S250, F40K modes) have been performed in a series of papers by Baranets et al. (2007, 2012, 2017) where the authors compared theoretically derived quantities with data obtained onboard IK-25 (VLF/HF emissions, electric and magnetic field fluctuations, plasma parameters) and remotely by MAGION-3 (HF emissions, accelerated electron fluxes, thermal plasma parameters). Presented results reflect the development of a beam or beam-anisotropic instabilities in the ionospheric plasma, some of them being a consequence of more complex combination (nonlinear) mechanisms of interaction, developing during injection of a beam into a beam.

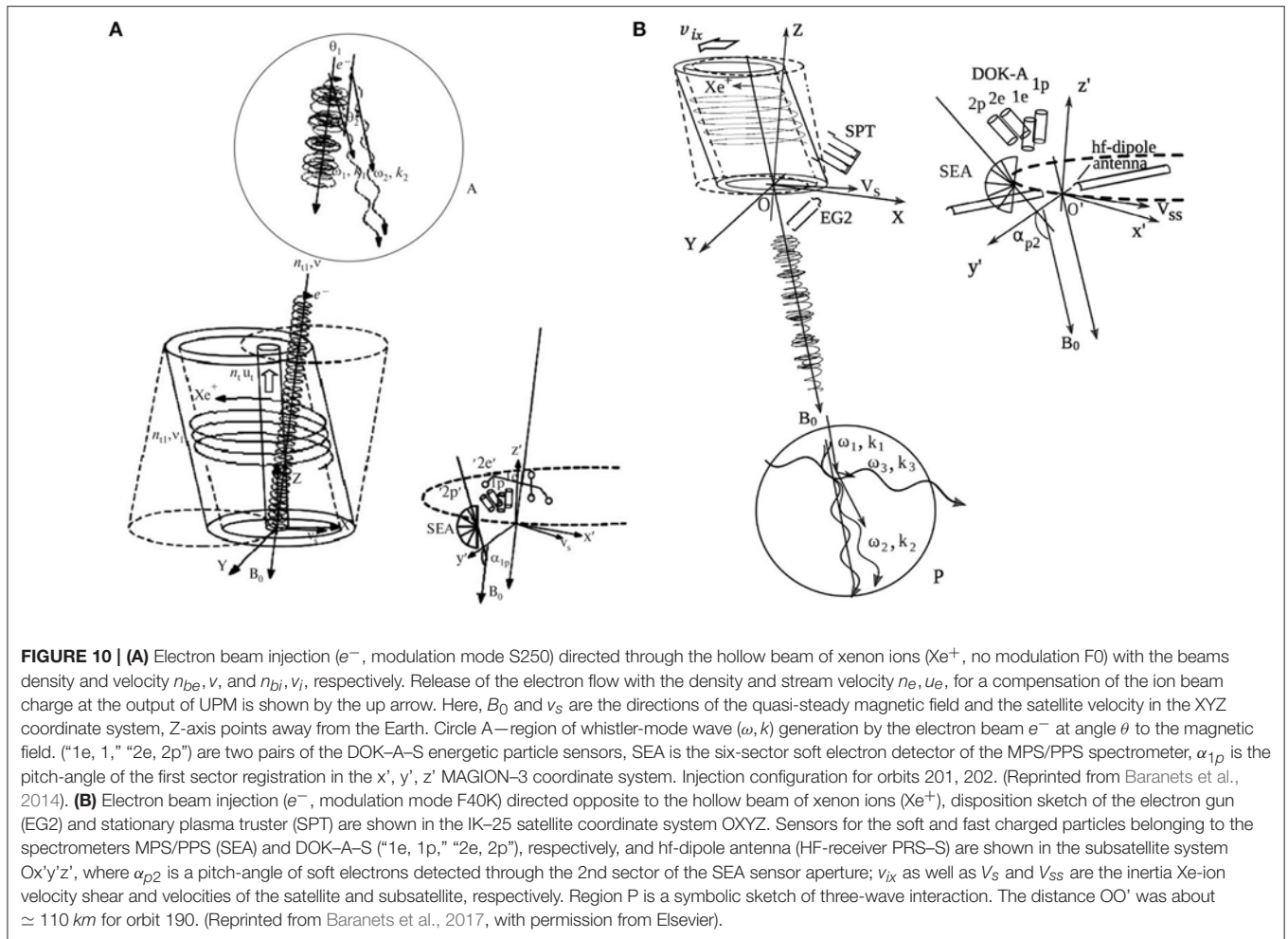
Baranets et al. (1999, 2000), and Oraevsky et al. (2001) analyzed waves generated during the DC and modulated electron beam injections in the near zone in a special dayside configuration when the electron beam was injected downward ($\alpha_{pe} \simeq 74 - 87^\circ$) and, asynchronously to the UEM cycle (S250 modulation mode), the neutral xenon/xenon plasma (F0 mode) were released upward ($\alpha_{pi} \simeq 121 - 132^\circ$), see also **Figure 3**. Beam-plasma discharge did not occur during this quasi-transverse injection, conditions for the beam-plasma instability development/suppression were investigated in these papers. Theoretical and numerical analysis of the observed data led the authors to the following conclusions:

- The beam plasma instability due to the electron beam injection ($\alpha_{pe} \simeq 74 - 87^\circ$) excited ELF/VLF waves activity below the ion cyclotron frequency, growth of the magnetic component of VLF waves was observed. Increase of thermal spread of the beam electrons leads to the suppression (decay) of the excited VLF waves.
- Magnetic disturbances were recorded, caused by either the beam primary electrons or return currents associated with the beam-plasma instability.
- The electron beam injection was accompanied by strongly anisotropic plasma heating and modulation of the plasma (ion) flows in the vicinity of the satellite.
- For an unmodulated electron injection, the efficiency of plasma (beam) heating with respect to the transverse component decreases substantially when $\alpha_{pe} + \Delta\alpha'/2 > 90^\circ$ ($\Delta\alpha'$ is the electron beam effective pitch-angle width).
- The thermal flux of O^+ and NO^+/O_2^+ ions decreased during the DC electron injections, probably as a consequence of the resonant coupling with the excited VLF waves.
- During the modulated beam-plasma interaction, the plasma resonance heating of T_{ex} , T_{ey} components at the modulation frequencies corresponding to the O^+ and NO^+/O_2^+ ion plasma frequencies were recorded.

Baranets et al. (2007, 2012, 2014) studied modulated electron injections (S250 mode) performed in the same direction (upward,

$\alpha_{pe} > 90^\circ$) as the xenon plasma release. Experimental observations of the anomalous fluxes of fast charged particles, disturbances of quasi-steady and ELF/VLF magnetic field components, and thermal plasma ion fluxes have been considered using the IK-25 and MAGION-3 data obtained from two consecutive orbits with similar ionospheric plasma and injection parameters. Anomalous large disturbances 400–500 nT and fluxes of fast electrons were registered by MAGION-3 (~ 100 km from IK-25) during these injections. In their analysis the authors did not consider the xenon plasma cloud as a diamagnetic cavity in the nearest vicinity of the spacecraft, describing it as a hollow beam of about half a kilometer diameter. The electron beam current profile was considered as a hollow flow after several gyro-turns due to electrostatic repulsion, electron beam modulation was commented in the papers. In such experimental configuration (**Figure 10A**) the system of the electron beam nested in the xenon ion beam was axially asymmetric with respect to the magnetic field direction due to the small velocity of xenon ions ($v_{iz}/u \sim 3 \times 10^4$), which was comparable to the velocity of the satellite ($v_{iz}/v_s \sim 1.5$) moving at an angle to the magnetic field. The beam of heavy xenon ions injected at the pitch angles range up to $\Delta\alpha_{pi} \simeq 60^\circ$ with a maximum flux density within the angles $\Delta\alpha_{pi} \simeq 30^\circ$ would play the role of a damping layer for waves induced by the electron beam in the entire interaction region in the vicinity of the satellite. Main attention was paid to study the electromagnetic and longitudinal waves excitation in different frequency ranges and the energetic electron fluxes disturbed due to wave-particle interaction with whistler waves. The authors described dispersion relations for the whistler-mode wave excitations for the beam-driven electromagnetic instability and estimated its growth rate (in a weak beam approach).

Recent papers (Baranets et al., 2012, 2017), analysing the stationary electron beam injection (F40K modulation mode, $2\mu s$ pulses) with simultaneous xenon neutral gas/plasma release (**Figure 10B**), reported a number of beam-plasma interaction effects both in the vicinity of the injecting satellite and in the far zone (subsattellite observations at distance about hundred kilometers, see also section 3.4.2). Based on thermal plasma ion fluxes, photometric, and wave data they found their correlation with beats of the density and velocity waves in the beam core in the high-frequency parallel and perpendicular wave fields. Theory and numeric estimations of correlation amplitudes are included. Onboard IK-25, sudden VLF amplification lasting ~ 7 – 8 s was recorded. Analysis shows the burst was observed near the linear stability boundary for the slow beam mode excited due to a dissipation of beam kinetic energy. HF waves recorded by MAGION-3 were disturbed during the electron beam/xenon injections. The amplification of the HF wave amplitude in the range $\sim 0.6 - 1.1$ MHz ($\omega_{ce}, \omega_{pe}, k_0 u$) and $2.0 - 2.3$ MHz ($\sim 2k_0 u$) were correlated with the VLF burst at IK-25. Narrow bursts within the $2.20 - 2.3$ MHz range were modulated by the VLF-LF wave activity in accord with the xenon ion beam current I_{bi} . Stimulated soft electron fluxes also recorded by MAGION-3 were explained as disturbed due to the Cherenkov wave-particle interaction (acceleration/scattering)— $50 - 70$ eV electrons interacted with excited waves near 0.1 MHz, fluxes



with energies $> 0.5 \text{ keV}$ could be stimulated by the beam-induced waves around $k_0 u$ near the linear stability boundary for longitudinal waves.

3.4. Distant Injection Effects

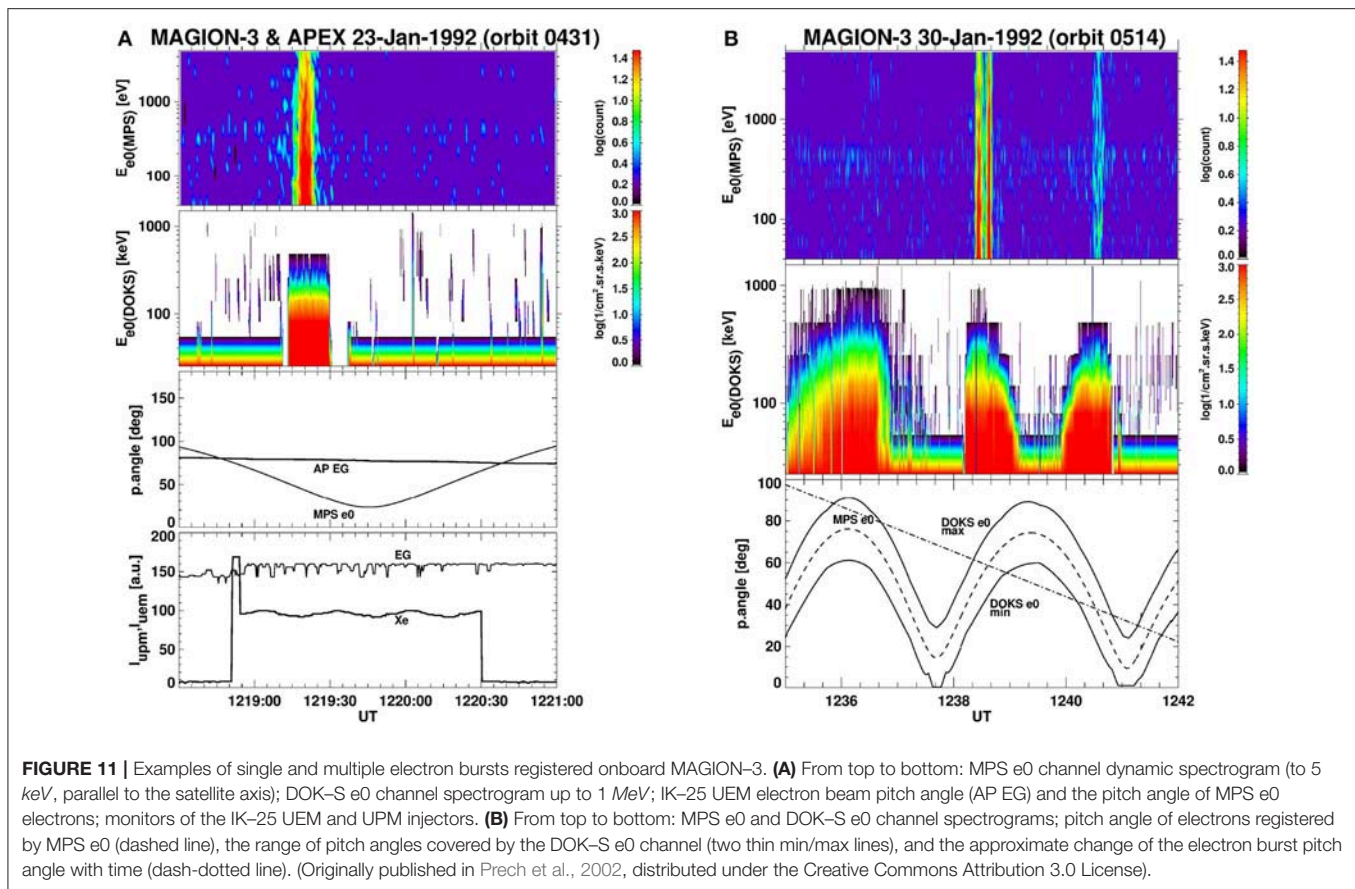
3.4.1. Energetic Particles

Nemecsek et al. (1996) and Prech et al. (2002) reported on observations of short intensive bursts of narrow-beam electrons onboard MAGION-3 during the electron injections from the main IK-25 satellite. The events were registered in the middle geomagnetic latitudes ($\text{INL} < 55^\circ$) near local noon when the two satellites moved approximately along the same magnetic meridian at relative distances 64–550 km. The localization of the burst observation to the altitudes from 800 to 1,500 km was probably a consequence of the local time favorable for the observation, the satellite orbit, and other parameters. The pitch angle distributions were narrow ($< 10^\circ$) and the electrons precipitated toward the atmosphere. The events lasted 4–10 s in the narrow MPS low-energy channel, but some events were doubled or tripled. The burst cross section was spatially limited to several tens of kilometers in the direction of the orbit and the electron energy extended to several hundred keV, without

measurable dispersion indicating proximity of the source. The events had different characteristics than electron microbursts already reported in the literature (for a short review see Prech et al., 2002). Examples of bursts registered by MAGION-3 are presented in Figure 11.

The main IK-25 satellite did not detect electron bursts of this kind as no other satellite orbiting in this region did but active electron injections from IK-25 were performed when the bursts were detected. The events of this kind were detected only during the UEM F40K mode operation and not during the S250 or other modulation mode injections. Neutral xenon or xenon plasma (UPM F0/F1000 modes) were simultaneously released so the quality of the satellite charge neutralization seems not to be important. Thirty four F40K-mode electron injections were scheduled on IK-25 during January and February 1992, but not all were 100% successful and with the MAGION-3 data coverage. The burst observation probability per session was about 0.25.

The energy-pitch angle properties of the bursts indicate according to Prech et al. (2002), that the bursts were created from low-temperature plasma or the beam emitted from the IK-25 satellite by some acceleration process effective only in one direction with respect to the magnetic field (e.g., field-aligned



potential drop), the spread of the pitch angles being caused by the original thermal velocity distribution. The authors were not in favor of the other possible source—ring current electrons released by some mechanism induced by particles emitted from IK-25. Questions if the beam-plasma interaction is able to accelerate the electrons to the observed energies, how the observed electrons moved across the magnetic field lines, and why the electron bursts were observed only in part of the active experiments are further discussed in the paper.

Baranets et al. (2012, 2014) reported on similar kind of fast electron fluxes (MAGION-3/DOK-S). The bursts had highly fluctuating wide (scattered) pitch-angle distribution, the electrons were simultaneously detected by both the DOK-S electron channels. The bursts were observed during the UEM S250 electron injection mode different from the previous paragraph case. According to the authors electromagnetic waves were excited via the electron-cyclotron resonance with the beam and subsequently scattered in a wave-particle interaction (Cherenkov resonance) giving rise to the energetic electrons distantly observed by MAGION-3.

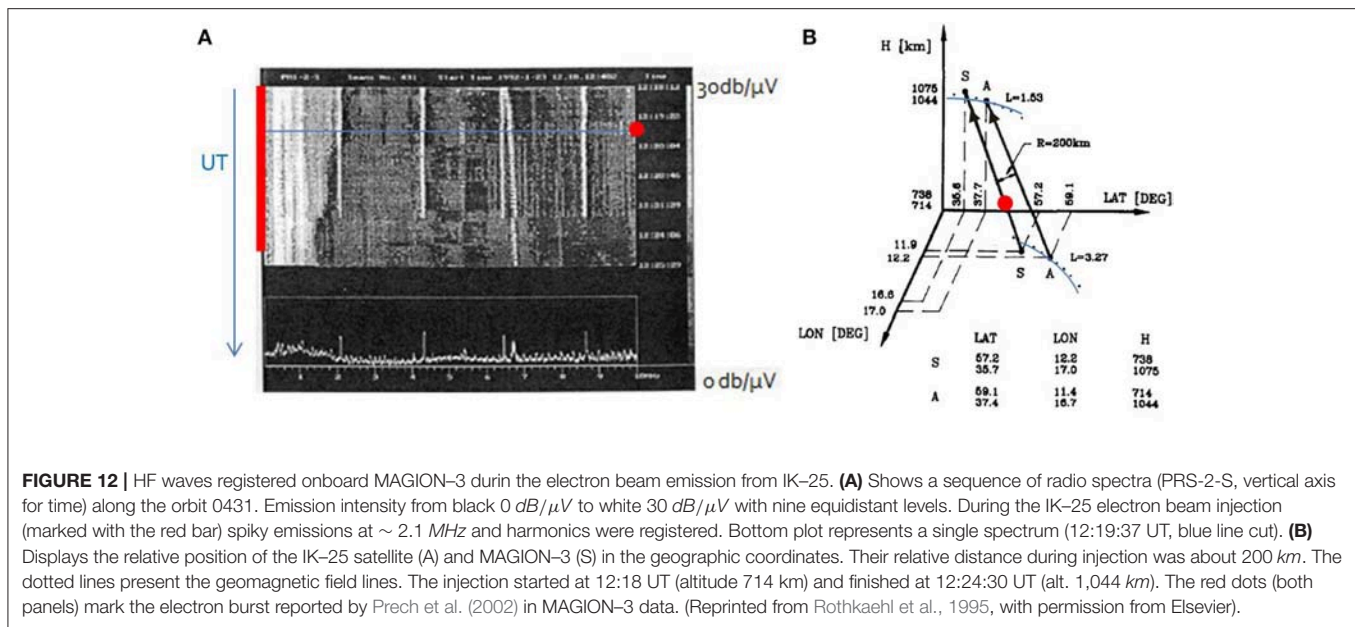
3.4.2. HF Waves

Rothkaehl et al. (1995) reported simultaneous observations of RF emissions on the mother IK-25 satellite and the MAGION-3 subsatellite during a downward UEM modulated electron beam injection (F40K mode, 2 μ s pulses, only neutral xenon

released by UPM during this case). The spacecraft separation was about 200 km. Strong electrostatic emissions at the upper-hybrid plasma frequency and its harmonics at the main spacecraft and spikes in the HF frequency range on MAGION-3 (see Figure 12) were simultaneously detected. The first emission peak was at frequency 2.15 MHz and three harmonics were also recorded. The authors concluded the electromagnetic emission of the oscillating current created by the injected electron beam was observed, the base frequency attributed to the electron beam current properties and the half-width of the emission line ~ 50 kHz referred to the electron gun pulse repetition frequency.

3.4.3. Ground VLF Transmitter Spectral Broadening

Spectral broadening of signals from ground-based VLF transmitters (16.5, 12.6, 11.333 kHz carriers) was observed onboard MAGION-3 during the modulated electron beam injection from the main IK-25 satellite (Oraevsky et al., 1994). The broadening of order of 300–500 Hz was apparently correlated with the 2-s cycle of the UEM gun pulses modulated in the range 30.5–31,250 Hz (2 μ s length) while no effects were observed for modulation frequencies 62.5, 125, and 250 kHz (Figure 13). The observations were made in the middle latitude ionosphere at altitudes 1,175–1,580 km, the distance between the two satellites was about 250 km. Scattering of whistler-mode waves into quasi-electrostatic waves by periodic small-scale plasma inhomogeneities or ELF plasma turbulence created by



the pulse-modulated electron beam were discussed as possible causes of the observed phenomenon. The authors compared the observation with previous rocket and space shuttle experiments and discussed both local (near the injection source) and remote (e.g., ionizing effect of the electron beam at the ionospheric E-region heights) mechanisms of plasma irregularities creation that had been suggested in literature.

3.5. Ionospheric Heating Experiments

Magnetosynchronous active experiments belong to a different group of active experiments in space plasma. One of possible approaches to investigate evolution of physical phenomena along a geomagnetic flux tube after a disturbance injection could be to use a special satellite orbit which runs in parallel with the geomagnetic field line for a sufficient time. These orbits are called magnetosynchronous and a concept of such active experiment in frame of the APEX programme was introduced by Ruzhin and Vaskov (1992). For the APEX orbit with inclination ~ 80° the magnetic meridian fell in the orbital plane about twice per day. The experiment was realized using the Dushanbe ionospheric heater for several IK-25 flyover orbits both near perigee and apogee.

Oraevsky et al. (1998b) reported on observation of the plasma barrier transparency effect (passing of a shortwave signal below the f_0F_2 frequency across the F2 layer, in the particular case for the heater frequency 5.98 MHz the barrier thickness was estimated ~ 100 km). The ballistic transport mechanism was suggested by the authors to explain detection of increased noise at frequencies around 6 MHz from PRS-3 data while the IK-25 satellite was magnetically connected to the heated part of the ionosphere. Oraevsky et al. (1998a) reported on similar results for the APEX and CORONAS satellites and the SURA ionospheric heating facility, more recently the experiment was repeated with the DEMETER satellite and the SURA heater (Zhang et al., 2016).

Variations of electron temperature and density near IK-25 apogee observed when magnetically connected to an ionospheric spot heated by the Dushanbe facility were reported and analyzed by Oraevsky et al. (1998b). Propagation of disturbances from locally artificially heated ionosphere into magnetosphere was modeled e.g., by Ruzhin and Vaskov (1992) and recently by Borisov et al. (2015).

4. COMPARISON WITH OTHER EXPERIMENTS AND GENERAL DISCUSSION

The analysis of the KM-10 floating probe potential and PEAS charge particle spectra has shown small charging and good neutralization during electron beam emissions from the IK-25 satellite. A larger spacecraft surface and methods of electron beam charge compensation made a difference comparing the electron beam emissions in a similar range of altitudes from the G-60-S sounding rocket (Managadze et al., 1988) or even at lower altitudes (SCEX-3, Mullen et al., 1991) where high body charging was observed.

Measurements of angular-energetic spectra of charged particles were rare prior to the APEX project. Previous sounding rockets and spacecraft active experiments included simple electrostatic analyzers that did not allow detail studies of pitch-angle distributions of charged particles in the vicinity of the beam emitting body. The space shuttle Spacelab-1/SEPAC and TSS-1/SETS active experiments (e.g., Burch, 1986; Watermann et al., 1988; Oberhardt et al., 1993) provided electron angular distributions with high time resolution but related pitch-angle distributions were not discussed. Moreover, configurations of these experiments were different from APEX. Most of the shuttle surface was covered by non-conducting ceramic tiles. The electron spectrometers were mounted on a platform inside

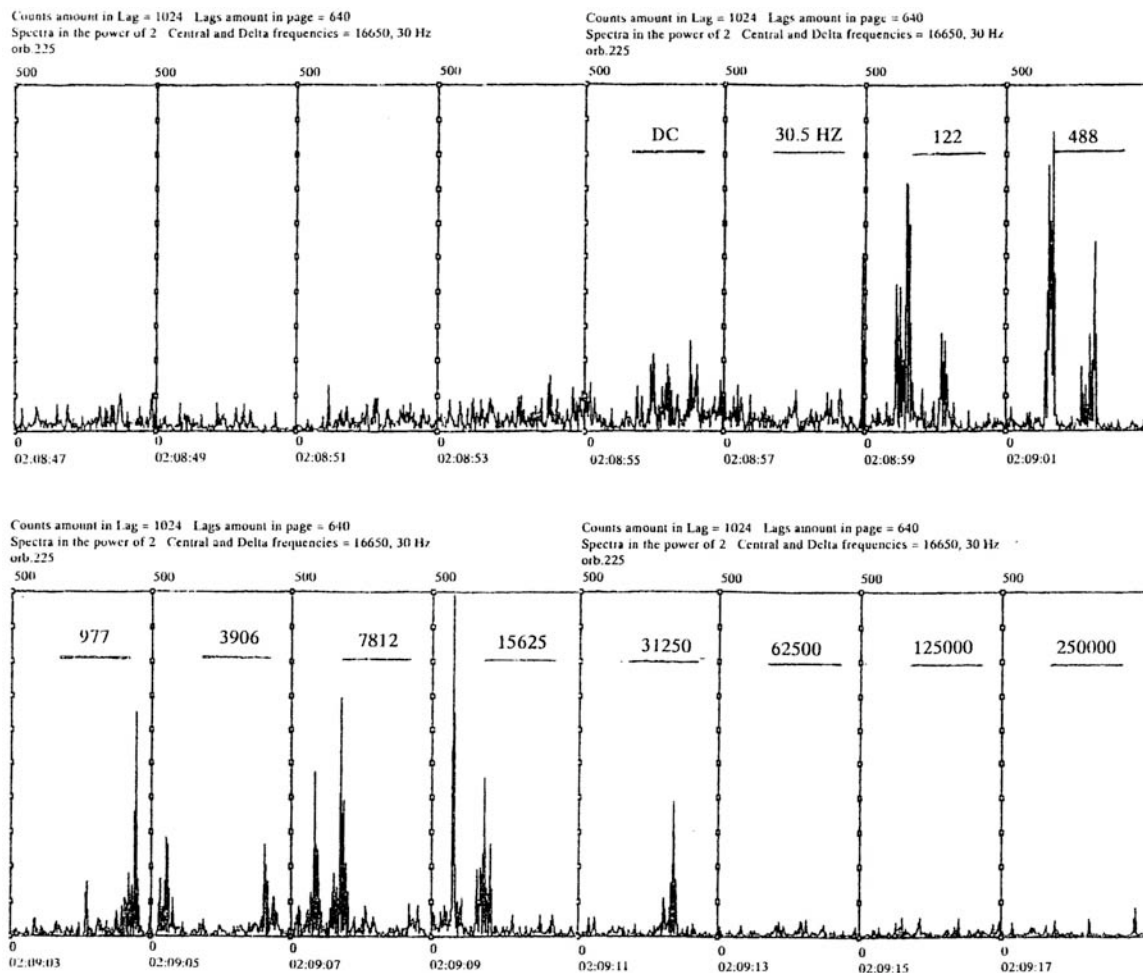


FIGURE 13 | Wave intensity registered by MAGION-3 at the frequency 16,650 Hz, i.e., 150 Hz away from the VLF transmitter carrier. The electron beam pulses are shown as horizontal bars with the modulation frequency indicated. (Reprinted from Oraevsky et al., 1994, with permission from Elsevier).

the shuttle payload bay where at least half of space was shielded. The angular distribution of the return current electrons during electron beam emissions did not show flux decrease in the shielded sectors indicating the origin of electrons near the platform. Their angular distribution depended more on the sector view angles against the platform normal and the direction toward the electron gun. For pitch-angles close to zero the electron energy was limited to about primary beam energy while for pitch-angles reaching $\sim 80^\circ$ intense electron flux was observed up to the spectrometer upper range (≈ 2.5 times the beam energy). In the APEX experiments, return electrons with energies of the primary beam and above were not observed in the IK-25 vicinity.

We consider important our observations of the short high-energy electron bursts in the remote zone by MAGION-3 during the electron injections. To our knowledge these observations have not been confirmed by other experiments yet, actually because no other experiment with two-point observations similar to APEX has been performed up to today. Comprehensive wave

and particle measurements were realized during the electron beam emissions from the STS-3 and Spacelab-2 space shuttle platforms (e.g., Gurnett et al., 1986; Banks et al., 1987; Bush et al., 1987) using the independent Plasma Diagnostic Package (PDP; mounted on a manipulator arm or free-flying, resp.) but they were made in the near zone at distance $\sim 10^1 - 10^2$ m. The MAGION-3 HF wave measurements also remain unconfirmed for the same reasons.

Kawashima and Akai (1986) reported on waves registered onboard the satellite JIKIKEN (EXOS-B) on an elongated orbit with apogee $\sim 25,000$ km that were excited during electron beam injections (much weaker than later used in the space shuttle or APEX experiments). They also observed waves near upper-hybrid or plasma frequencies, electron-cyclotron frequency, and LF frequencies related to the electron beam instrumental modulation. Accompanying harmonics were supposed to be generated instrumentally due to the saturation of the signal level. Harmonics of the plasma frequency in the APEX data might have the same instrumental origin which was not discussed by Budko

et al. (2003), or they could truly exist due to the beam bunching as noted by Kiraga et al. (1998).

From their analysis of the return electron spectra, Nemecek et al. (1997) and Budko et al. (2003) deduced a presence of electric fields of a strength up to 100 V/m in the vicinity of the spacecraft that was sometimes questioned by a part of the scientific community and would require further study. Anyway, according to the 3D computer simulation by Pritchett (1991) electric fields of the order of tens of volts per meter can be expected during an electron beam injection in similar conditions.

Complexity of the IK-25 design (structure, surfaces, etc.) and a lack of related engineering data have not allowed to synthesize a more realistic model of the electron beam and xenon plasma injections and return currents in the near zone including the satellite influence that would help to interpret the observational data. Earlier papers related to the APEX wave observations (Kiraga et al., 1998; Budko et al., 2003; Kiraga, 2003) did not aim to establish a comprehensive theory of wave generation by the APEX electron beam and xenon plasma injections but in their analysis of observed data they rather referred to previous theoretical treatments of beam-plasma interaction which is non-linear in the APEX case. Full understanding of the coexistence and competition between various instabilities requires dedicated simulations which would take into account many unique characteristics of the APEX experimental setup, but such APEX-related simulation models have not been developed. Baranets et al. (2007, 2012, 2017) presented non-linear theoretical analysis of beam-plasma interaction for systems of electron beam nested in or opposite to the ion beam that were related to the APEX injection configurations. Fast electrons from the MAGION-3 observations were explained as a result wave-particle interaction: whistler waves excitation by the electron-cyclotron resonance and subsequent scattering. To verify their results, a detail knowledge on wave propagation (k vector) would be useful, but such data were not gathered in the APEX project.

The APEX project suffered from a relatively short life-time of active devices UEM-2 and UPM. The active experiments with charge particle beam injections were performed only during the first six months after the launch, recurrent technical issues precluded to obtain complementary data during the second half year or to repeat unsuccessful active experiment configurations. Above that, some operational and science information was not available (fast beam current monitoring, optical diagnostics, cold plasma spectra), the field and wave instruments did not provide complete information. Technical difficulties connected with the precise control of the MAGION-3 orbit did not allow intended near- and mid-zone observations during the active experiments while the hundred-kilometers satellite distances were preferred for two-point passive measurements in the auroral region. Last but not least, the available computer network communication means limited daily supervision by scientists, quick preliminary data analysis, and feedback toward the main satellite operation

control—reasonable demands of a contemporary active experiment.

Many questions connected to the APEX active injections have not been answered, e.g., what was the size and structure of the electron acceleration region or what conditions allowed the observed acceleration up to relativistic energies. A new mission could provide multi-point measurements in medium and distant zones of injection ($1-10^3$ km) using a constellation of nanosatellite probes that would allow the study of waves growth and propagation and electron acceleration simultaneously at different distances along and perpendicular to the beam. If feasible, the orbit (active experiment parts) should be chosen so as the emission pitch-angle, altitude, latitude, L shell, and local time are not closely bounded, with sufficient coverage of parameter space and repeatability of emission parameter subsets. Wave-particle correlators and ion analyzer with mass selection would be a useful extension of the APEX scientific payload. A new experiment should be also equipped with a 3D measurement of LF and HF electric and magnetic fields. Such data, when processed using the latest analysis methods to obtain wave propagation direction (e.g., Santolik et al., 2003), polarization (Santolik et al., 2016; Taubenschuss and Santolik, 2018), and HF mode identification (Santolik and Parrot, 2006), can give a new insight into the beam-plasma interaction processes. Last but not least, Kiraga (2003) discussed also usage of the APEX type electron injection and HF waves analysis for monitoring density of cold plasma.

5. CONCLUSIONS

The broadly focused APEX active experiments have an important place in the long line of active space experiments. They brought new results or they complemented results of previous projects made using sounding rockets, space shuttles, or scientific/technological satellites in ionosphere and magnetosphere. The main achievements of the active experiments within the APEX project could be listed as

- successful neutralization of the spacecraft charge during electron beam emissions by xenon/xenon plasma releases at altitudes 500 – 1,000 km,
- analysis of ELF/VLF and HF wave activity near the electron-beam emitting satellite,
- first *in-situ* measurements in the distant zone of injection (MAGION-3 subsatellite),
- MAGION-3 observation of electron bursts accelerated to energies of several hundred keV during the electron beam injections from the main satellite.

The mother—daughter (multi-point) satellite projects were quite rare till the end of the twentieth century and we can treat the APEX project also as the pathfinder in this methodology and the MAGION series satellites as predecessors of contemporary micro, nano, and cube satellites. Repeating the APEX project with state-of-the-art scientific

equipment and a fleet of small satellite probes around would certainly bring new scientific achievements and help to answer some still open questions in the field of space plasma.

AUTHOR CONTRIBUTIONS

YR, VD, ZN, and JS contributed to the conception and design of the APEX project. LP wrote the first draft of the manuscript. YR wrote substantial parts of several sections of the manuscript. LP, JS and ZN selected and prepared figures to be reprinted in this review. All authors contributed to

the manuscript revision, read, and approved the submitted version.

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Recalling and Updating Research on Diamagnetic Cavities: Experiments, Theory, Simulations

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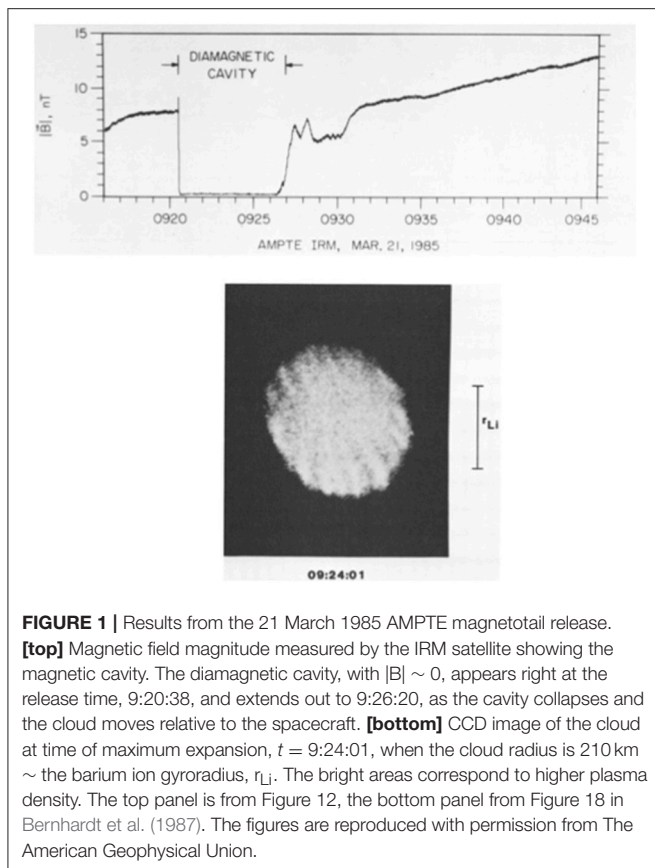
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In the decade from the mid 80's to the mid 90's there was considerable interest in the generation of diamagnetic cavities produced by the sub-Alfvénic expansion of heavy ions across a background magnetic field. Examples included the AMPTE and CRRES barium releases in the magnetotail and magnetosphere as well as laser experiments at various laboratories in the United States and the Soviet Union. In all of these experiments field-aligned striations and other small-scale structures were produced as the cavities formed. Local and non-local linear theory as well as full particle (PIC), hybrid, and Hall-MHD simulations (mostly 2-D) were developed and used to understand at least qualitatively the features of these experiments. Much of this review is a summary of this work, with the addition of some new 3-D PIC and Hall-MHD simulations that clarify old issues associated with the origin and evolution of cavities and their surface features. In the last part of this review we discuss recent extensions of the earlier efforts: new space observations of cavity-like structures as well as new laboratory experiments and calculations with greatly improved diagnostics of cavities formed by expansions of laser-produced ions at super-Alfvénic speeds both across and along the background magnetic field.

Keywords: magnetic cavities, plasma instabilities, active experiments in space, kinetic plasma simulations, Hall-MHD simulations

INTRODUCTION

Many active experiments in space involve the release of canisters of neutral barium atoms. Barium has the very interesting property that it has a very long photo-ionization time, ~ 30 s. This allows the formation of large clouds of barium ions expanding across the ambient magnetic field. Such ion expansions exclude the magnetic field, creating a magnetic cavity. The Active Magnetospheric Particle Tracer Experiment (AMPTE) mission launched in 1984 (Valenzuela et al., 1986) and a second mission, the Combined Release and Radiation Effects Satellite (CRRES) launched in 1991, which released smaller amounts of chemicals, were carried out to provide information of how newborn ions in the magnetotail or upstream of the bow shock traveled throughout the magnetosphere. For example, AMPTE produced artificial comets with the release of barium as well as lithium upstream and just behind the bow shock. Similarly, there were releases in the magnetotail; the barium ions produced large clouds (radius ~ 200 km) that were visible from the Earth. The top panel in **Figure 1** shows the magnetic field magnitude as a function of time as measured by the magnetometer on the release module (Lühr et al., 1988) during the first magnetotail experiment on March 21, 1985 (Figure 12, Bernhardt et al., 1987). The diamagnetic



cavity, with $|B| \approx 0$, appears shortly after the release time, 9:20:38, and extends out to 9:26:20, as the cavity collapses and the cloud moves relative to the spacecraft. Optical measurements from the ground confirm the density and velocity profiles of the barium as a function of time. Cavity Formation and Properties of this review provides a simple derivation of the dynamics of the cavity formation that has been shown to be consistent with the AMPTE data. Results from experiments of cavity formation performed in the laboratory using a source of expanding energetic ions from a target irradiated by a laser will also be discussed that suggest the commonality of the diamagnetic cavity formation process.

The lower panel of **Figure 1** is a photograph of the same AMPTE barium cloud obtained from a CCD camera on the ground in White Sands, NM (Figure 18, Bernhardt et al., 1987). This photo was taken at about the time that the cloud has reached maximum expansion across the magnetic field and shows field-aligned striations on the surface of the cloud. Striations are commonly observed on barium clouds, but most often seen at later time at longer wavelengths when the cloud collapses radially and elongates along the ambient magnetic field. Short wavelength striations are also frequently observed in laboratory experiments. The origin and properties of these structures are the subject of Striation Growth and Evolution.

It is interesting to note that historically the chemical release experiments in space (AMPTE to CRRES) occurred in about the same time frame as the laboratory experiments were being fielded in both the U.S. and the Soviet Union. The experiments will

be described in Cavity Formation and Properties and Striation Growth and Evolution, and have been discussed in reviews by Akimoto et al. (1988), Zakharov (2003), and Zakharov et al. (2006). These two groups of activities were not programmatically related, and much theory and simulation work was directed toward one or the other type of experiment. As we show later, the theory and simulation program at the Naval Research Laboratory (NRL) was active in both the AMPTE and CRRES missions, and was also the home of a significant laser experimental effort, so they were well-situated to connect these two seemingly different efforts.

In contrast to this experience track is one recently proposed by Howes (2018). In his review he argues that laboratory experiments can now replace space experiments, because laboratory experiments are cheaper, are easier to field, have more controllable boundary and initial conditions, and are more readily reproducible/repeatable. He gives examples of laboratory experiments that are addressing issues of direct importance to space plasmas physics. These include: plasma turbulence, magnetic reconnection, particle acceleration, collisionless shocks, kinetic, and fluid instabilities, as well as other processes. He does recognize, however, that space experiments do have larger spatial and longer temporal domains, use smaller size probes relative to plasma scales and provide 3-D velocity diagnostics that are important for some types of experiments. In this review, we discuss the formation and evolution of diamagnetic cavities—a very active area of research three decades ago in which there were significant and complementary efforts in space and laboratory experiments. Interest does continue at present, albeit in somewhat different physical regimes in the laboratory, but without corresponding active space experiments. While the Howes review (2018) discusses how laboratory experiments can now be used to understand space plasma physics issues, the present paper stresses the important role of the interplay of space and laboratory activities that led to basic understanding of diamagnetic cavities in the previous generation. As we emphasize in this review, a broad-based array of both laboratory and space experiments addressed by a wider range of theoretical and simulation techniques that can span both spatial and temporal scales of laboratory and space conditions could be beneficial to the needs of different funding sponsors.

The fundamental question we address in this review is what we did, or did not, know concerning the nature of diamagnetic cavities by 1993, and how recent work has improved this understanding. The basic goals of this review related to this overarching question are as follows. First, we recall the basic physics of cavity formation and look back at the original experiments that validated this model (Cavity Formation and Properties). Second, we review a variety of previous experiments conducted in both space and in the laboratory to investigate the main dynamical features of cavities—namely the formation of field-aligned surface striations. We discuss related theory and simulations that describe at least qualitatively the mechanisms for the growth of these surface waves. And we address recent experiments and 3-D simulations that have improved this overall understanding (Striation Growth and Evolution). Third, we discuss present-day interest in cavities—new phenomena in

space and laboratory experiments and simulations to investigate magnetic cavities in different regimes (Extensions). Finally, in Conclusions we provide a short summary and suggest where research in diamagnetic cavity phenomena may be headed in the future.

CAVITY FORMATION AND PROPERTIES

The formation of a diamagnetic cavity, such as the AMPTE magnetotail magnetic cavity shown in **Figure 1**, can be understood from a simple pressure balance argument based on conservation of energy (e.g., Gisler and Lemons, 1989; Ripin et al., 1993). We consider the 3-D radial expansion of an ideal, perfectly conducting dense plasma of radius R , velocity V and total mass M in a uniform magnetic field B_0 . The properties of the expansion are determined by the work the plasma does in pushing the magnetic field out of the volume that it occupies. The final size of the cavity results from the equilibration between the (thermal) pressure of the expanded plasma and the external magnetic field pressure. The stopping radius (R_B) and the duration of the expansion (t_0) are found using the conservation of energy E_0 :

$$E_0 = \frac{1}{2}MV^2 + \frac{B_0^2}{2\mu_0} \frac{4\pi R^3}{3}, \quad (1)$$

where initially $V(t=0) = V_d$, $R(0) = 0$ and thus $E_0 = MV_d^2/2$. The expansion continues until $t = t_0$, when $V(t_0) = 0$, $R(t_0) = R_B$, with $E_0 = (B_0^2/2\mu_0)(4\pi R_B^3/3)$. From equating these two expressions for E_0 , we find $R_B = (3\mu_0 MV_d^2/4\pi B_0^2)^{1/3}$; R_B is referred to as the magnetic confinement radius. In this simple picture, we neglect any additional thermal energy or internal Ohmic heating. In the next section we will go beyond this model to look in detail at the current on the edge of the expanding plasma and how it leads to cavity formation and plasma heating.

Later, we will also need to know the rate of deceleration of the plasma cloud. By taking the time derivative of Equation (1), with $dE_0/dt = 0$, we obtain

$$g = -\frac{dV}{dt} = \frac{2\pi B_0^2}{\mu_0 M} R^2 = \frac{3}{2} \left(\frac{R}{R_B} \right)^2 \frac{V_d^2}{R_B} \quad (2)$$

(with $g > 0$), which evidently increases as R^2 and reaches its maximum at $R = R_B$,

$$g_{\max} = \frac{3}{2} \frac{V_d^2}{R_B}. \quad (3)$$

One can also integrate to get $R(t)$ and $V(t)$ and find $t_0 \approx 1.3 R_B/V_d$.

In similar fashion in the case where the expansion is cylindrical in two dimensions (with cylinder of length L), so that the excluded magnetic energy is $B_0^2 r^2 L \pi / 2\mu_0$, we can show, with $R_B = (\mu_0 M V_d^2 / B_0^2 L \pi)^{1/2}$, that

$$g = -\frac{dV}{dt} = \frac{B_0^2 \pi L}{\mu_0 M} R = \left(\frac{R}{R_B} \right) \frac{V_d^2}{R_B}, \quad (4)$$

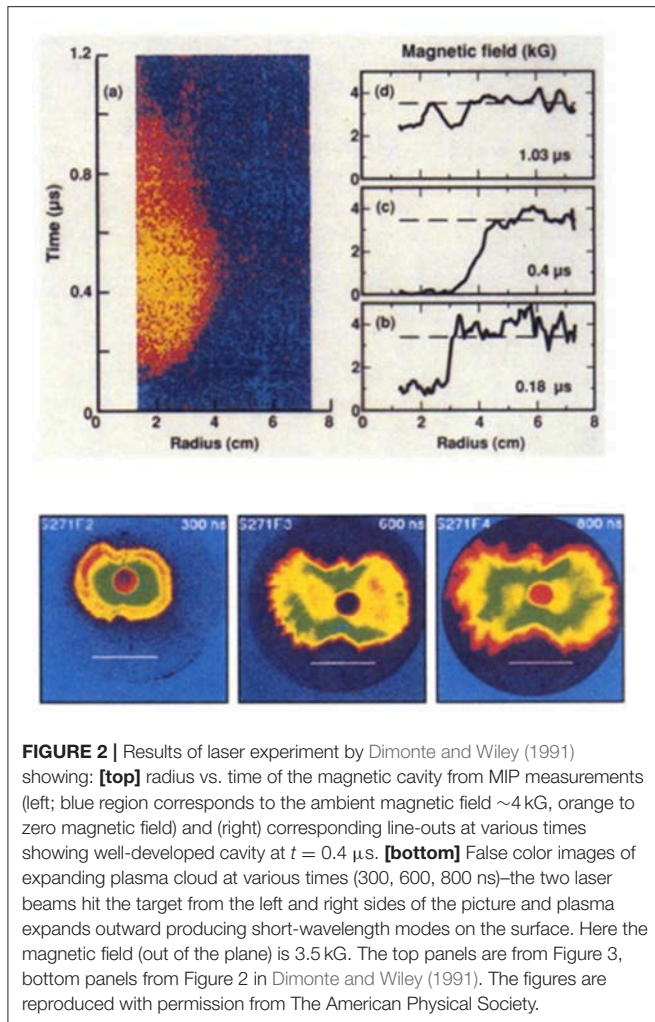
so that $g_{\max} = V_d^2/R_B$ at $R = R_B$ and $V_d t_0/R_B = \pi/2$. The difference in the rate of deceleration of the plasma near the end of the expansion between 2-D ($\sim R$) and 3-D ($\sim R^2$), and the maximum deceleration $g_{\max}/(V_d^2/R_B) = 1$ in 2-D and 1.5 in 3-D, will appear in the development of the surface striations is shown in the next section. It should also be noted that in 2-D a slightly different result occurs if free expansion is allowed in the axial direction (Gisler and Lemons, 1989). Also at least in simulation, it is possible to have a one-dimensional expansion as well, with the excluded magnetic energy, $B_0^2 L^2 R / 2\mu_0$. In this case the velocity decreases linearly in time, $V_d t_0/R_B = 0.5$ and $g = 0.5 V_d^2/R_B$ is a constant.

In the AMPTE case, ground based observations indicated that the maximum cavity radius was about 210 km. Estimating the initial plasma mass is a bit difficult because only a fraction of the barium atoms that are released become ionized. Ground-based measurements could also confirm that the expanding barium plasma compressed into a thin shell, leading to a narrow current layer. This current loop provides the magnetic field that cancels out the background field inside the loop, forming the diamagnetic cavity. Like the Lühr et al. (1988) observations of the AMPTE diamagnetic cavity measured by the release spacecraft (IRM) shown in **Figure 1**, direct spacecraft measurements of the diamagnetic cavity formation from smaller barium releases in the CRRES program were also carried out (e.g., Bernhardt, 1992).

Diamagnetic cavities were also produced in laboratory experiments in about the same time frame as the AMPTE barium releases. For example, at the Naval Research Laboratory the 30 J Pharos III glass laser was used to illuminate an aluminum target in a low density plasma embedded in a strong magnetic field ($B \sim 0.1$ – 1 kG) (Ripin et al., 1987, 1990, 1993). Magnetic probes were used to measure the magnetic field excluded by the expanding target plasma, although the initial measurements were not very accurate. Another experiment (Dimonte and Wiley, 1991) employed the two-beam 200 J Janus laser and various target materials at Lawrence Livermore National Laboratory and included a magneto-optic imaging probe (MIP) that was developed and used to obtain accurate measurements of the magnetic field. This technique uses Faraday rotation to measure the magnetic field profile continuously in space and time, as shown in the top panels of **Figure 2** (Figure 3 in Dimonte and Wiley, 1991). The top left panels present streak camera images across a transverse cut vs. time, showing the cavity expanding to a maximum radius (~ 3 cm) and then collapsing. The line-outs at various times in the right side of the top panel show that the field is indeed totally excluded in the cavity and the magnetic field gradient at the cavity boundary is relatively steep. Other experiments were carried out in Novosibirsk, USSR, using a 1 kJ CO₂ laser and observed similar behavior (Zakharov et al., 1999; Zakharov, 2002).

STRIATION GROWTH AND EVOLUTION

While the formation of a magnetic cavity is consistent with conservation of energy and magnetic field exclusion in both space and laboratory experiments was expected, the development



of field-aligned flute modes on the surface of the expanding plasma cloud was a more surprising feature. One example is the CCD image of the AMPTE magnetotail release that was already shown in **Figure 1**. With *in situ* measurements from a single spacecraft it would have been difficult to infer the character of these structures. In this section we discuss the experiments, theory and simulations, which led to a good understanding of how such surface features develop, in the years during and immediately after the space chemical releases and the laboratory laser experiments were carried out. We conclude with a discussion of more recent high-resolution experiments and simulations that have further enhanced our knowledge of the underlying processes.

Experiments

In some laser experiments small-scale short-wavelength striations appeared at the edge of the cavity, reminiscent of the AMPTE release shown earlier in **Figure 1**. Examples include the early small-scale experiments by Okada et al. (1981, Figure 3) the Ripin et al. experiments (Figure 6, 1993), Zakharov et al., experiments (Figure 3, 1996), and also the Dimonte and Wiley

(1991, Figure 2) experiments. An example of this type of behavior from the Dimonte and Wiley experiment is shown in the bottom panel of **Figure 2** displaying false color images of expanding plasma cloud at various times (300, 600, 800 ns). In this experiment the two laser beams hit the target from the left and right sides of the picture and plasma expands outward producing short-wavelength modes on the surface. However, in other experiments, larger flutes were observed (e.g., Ripin et al., 1990 Figure 2; Ripin et al., 1993, Figure 3), which showed unusual, finger-like projections. Similarly, in some experiments where large flutes were observed, the overall cavity size was smaller (i.e., $r < R_B$) and the flutes extended some distance beyond $r = R_B$. This suggested that another experimental parameter also plays an important role. As originally shown by Zakharov et al. (1986), this parameter is the ratio of the gyroradius of the expanding target ions (of charge $Z_d e$ and mass m_d), $\rho_d = V_d / \omega_{cd}$ (where the ion gyrofrequency is $\omega_{cd} = Z_d e B_0 / m_d$) to R_B . For small ρ_d / R_B , small-scale flute modes are generated and cavity sizes are $R \sim R_B$; larger flutes and smaller cavities result when $\rho_d / R_B \sim 1$. Dimonte and Wiley (1991) changed the magnetic field, the laser energy and the target material in their experiment to vary ρ_d / R_B and modify the size of the resulting cavities to verify these predictions. The ratio ρ_d / R_B determines how magnetized are the target ions during the expansion, which can be expressed in terms of the size of the Hall-term in a Hall-MHD description of the cavity dynamics, as discussed later. In cases of very energetic target ions, where $\rho_d / R_B \gg 1$, the expanding plasma generates a jet-like penetration across the magnetic field and creates only a small magnetic cavity, as shown in experiments by Mostovych et al. (1989) and Plechaty et al. (2013). Short-wavelength surface modes that are observed in this regime are likely due to a lower hybrid velocity shear (i.e., Kelvin-Helmholtz-like) instability (Peyser et al., 1992).

Further analysis of the flutes observed in the NRL experiments suggest that they could appear before the cavity reached its final size, and there could be a turn-on condition that in some experiments seemed to be independent of the magnetic field (Ripin et al., 1993). There were also suggestions that as the flutes grew, their wavelengths increased, although it was not clear whether this occurs by mode coupling, mode coalescence, or some other non-linear process. Dimonte and Wiley (1991) were able to obtain wavenumber spectra that showed such a shift to longer wavelengths. Other observations showed that flutes seem to bend in the direction of electro gyro-motion and the tips of the flutes were observed at times to bifurcate (Ripin et al., 1993). Again, we emphasize that the reproducibility of these laboratory experiments and the ability to vary parameters independently were important to document the behavior of the surface striations as they grew and developed, perhaps non-linearly.

Theory

In the pressure-balance model presented in Cavity Formation and Properties, the expanding plasma slows down as it does work in excluding the background magnetic field to form a diamagnetic cavity. The deceleration of the spherically expanding plasma ($-g, g > 0$) is given by Equation (2). The plasma is also heated, and equilibrium is achieved when the thermalized sphere

of plasma is in pressure balance with the external magnetic field. Here we go beyond the simple dynamical model and look in more detail at the physics of the expanding plasma to understand the generation of the plasma current that leads to the expulsion of the magnetic field, the formation of flute modes and the resulting thermalization of the plasma.

At the leading edge of the expanding plasma cloud there is a density gradient ε_n given by

$$\varepsilon_n = -\frac{1}{n(r)} \frac{dn(r)}{dr}, \quad (5)$$

defined so that $\varepsilon_n > 0$ at this interface. One then expects that the interface between the expanding but decelerating plasma and the magnetic field would be unstable to a Rayleigh-Taylor instability, the growth rate of which according to MHD is

$$\gamma_{RT} = (g\varepsilon_n)^{1/2}. \quad (6)$$

However, in the case of expanding cavities produced in active experiments, the cavity is usually small compared to the wavelength of the Rayleigh-Taylor instability and the observed waves on the surface of the cavity are evidently smaller than the cavity diameter. In this case one needs to include the Hall term in the wave analysis. With the addition of this term, Hassam and Huba (1987) and Huba et al. (1987) have shown that a shorter wavelength, faster growing instability occurs with growth rate

$$\gamma_{LLR} = k(g/\varepsilon_n)^{1/2}. \quad (7)$$

In contrast to the usual Rayleigh-Taylor instability, in this “Large-Larmor Radius” (LLR) Rayleigh-Taylor instability the character of the unstable modes is different: the fluid motion of the plasma is current-free, rather than divergence-free, producing finger-like projections of the plasma into the magnetic field region, rather than the usual spike-and-bubble configuration characteristic of a Rayleigh-Taylor instability. **Figure 3** shows the growth rates of both the Rayleigh-Taylor and the LLR instability as a function of the inverse density scale length ε_n . These solutions are derived from solving a set of single fluid MHD equations (blue curve) as well as equations containing the Hall term (red curve) (Hassam and Huba, 1987).

An alternative way to understand this new type of Rayleigh-Taylor-like instability is from a kinetic point of view (Winske, 1988, 1989). The expanding plasma compresses into a thin shell as the ions exclude the magnetic field and a radial electric field E_r (< 0) drags the electrons and magnetic field along to keep the plasma quasi-neutral, as shown in the left panels **Figure 4** (Figure 1 in Winske, 1989). In terms of the radial ion momentum equation (for unmagnetized ions),

$$n_i m_i \frac{dV_{ir}}{dt} = n_i e Z_d E_r - \frac{dP_i}{dr}, \quad (8)$$

with the ion pressure given in terms of the ion temperature, $P_i = n_i(r)T_i$ with $T_i = m_i v_i^2/2$ assumed simply as a constant, the radial

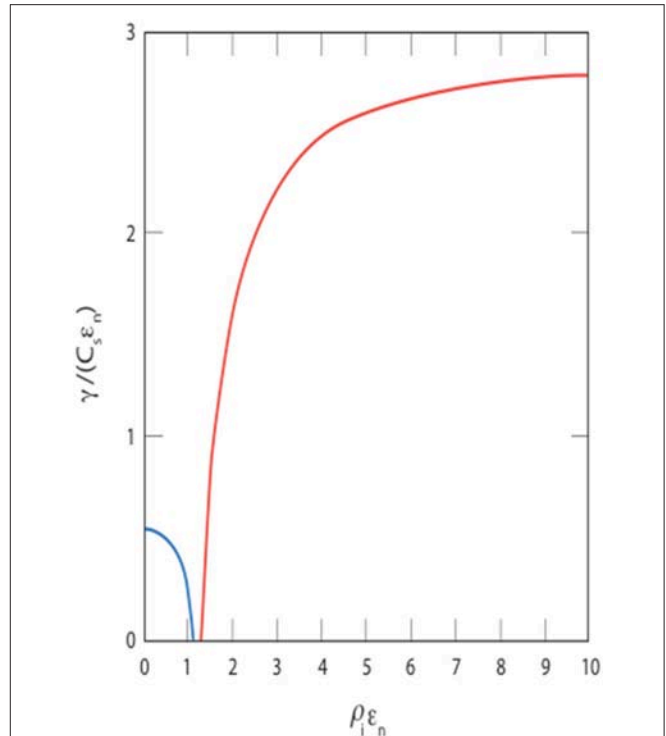


FIGURE 3 | Calculated growth rate of the usual Rayleigh-Taylor instability (blue curve) and the Large Larmor Radius Rayleigh Taylor instability (red curve) as a function of $\rho_i \varepsilon_n$. Solutions come from solving set of single fluid MHD Equations (blue curve) with the addition of the Hall term (red curve). The figure is redrawn from the original Figure 1 in Hassam and Huba (1987), with permission from The American Geophysical Union.

electric field can be written as

$$E_r = \frac{T_i}{Z_d e n_i} \frac{dn_i}{dr} + \frac{m_i}{Z_d e} \frac{dV_{ir}}{dt} = -\frac{m_i v_i^2}{2Z_d e} \varepsilon_n - \frac{m_i}{Z_d e} g. \quad (9)$$

In the rest frame of the ions, the electrons ($\mathbf{E} \times \mathbf{B}$) drift in the azimuthal direction

$$V_{ExB} = -\frac{E_r}{B_0} = \frac{g}{\omega_{cd}} + \frac{1}{2} \frac{v_i^2 \varepsilon_n}{\omega_{cd}} = V_g + V_n, \quad (10)$$

due to a combination of the “gravitational drift” V_g and the ion diamagnetic drift V_n . This azimuthal, relative electron-ion drift provides the current that produces a magnetic field which opposes the background magnetic field leading to the formation of the cavity. The electron-ion drift also gives rise to a fast-growing electrostatic instability. This instability was originally discussed by Davidson and Gladd (1975) and Gladd (1976), who studied decelerating plasma sheaths ($g > 0$) in imploding theta-pinch plasmas, and Okada et al. (1979), who suggested that the instability could occur in expanding magnetized laser plasmas. The geometry for the linear analysis is shown in the right side of **Figure 4** (Figure 1 in Winske, 1988). With $V_g = 0$, the resulting instability is the well-known lower hybrid drift instability (Krall and Liewer, 1971); the real frequency ω_r , the maximum growth

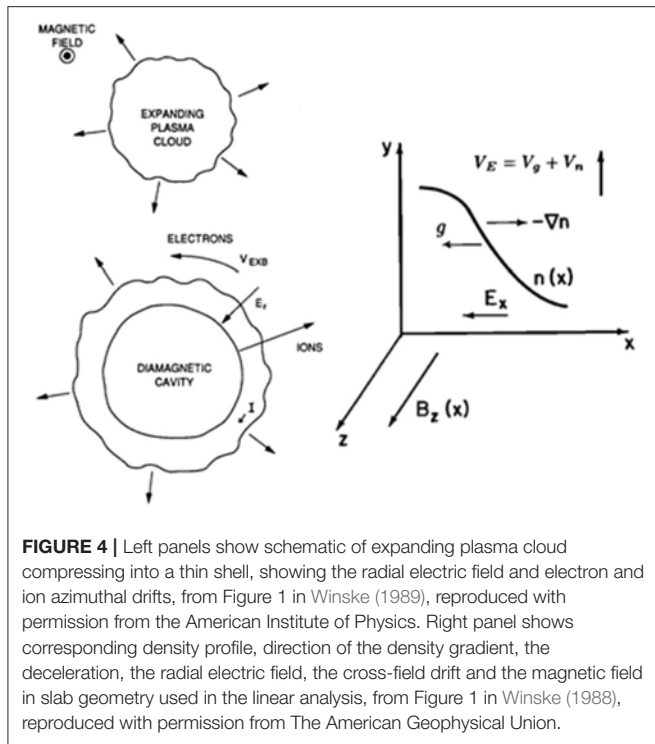


FIGURE 4 | Left panels show schematic of expanding plasma cloud compressing into a thin shell, showing the radial electric field and electron and ion azimuthal drifts, from Figure 1 in Winske (1989), reproduced with permission from the American Institute of Physics. Right panel shows corresponding density profile, direction of the density gradient, the deceleration, the radial electric field, the cross-field drift and the magnetic field in slab geometry used in the linear analysis, from Figure 1 in Winske (1988), reproduced with permission from The American Geophysical Union.

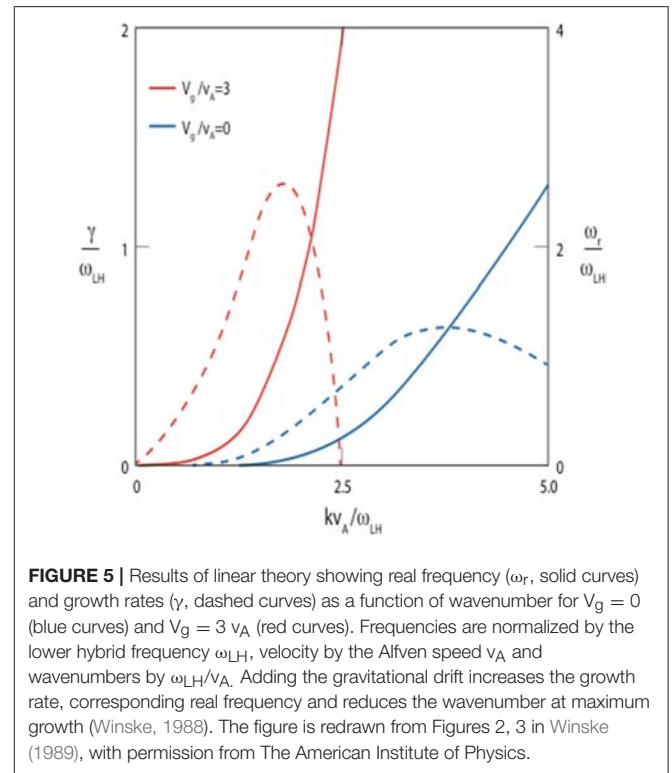


FIGURE 5 | Results of linear theory showing real frequency (ω_r , solid curves) and growth rates (γ , dashed curves) as a function of wavenumber for $V_g = 0$ (blue curves) and $V_g = 3 v_A$ (red curves). Frequencies are normalized by the lower hybrid frequency ω_{LH} , velocity by the Alfvén speed v_A and wavenumbers by ω_{LH}/v_A . Adding the gravitational drift increases the growth rate, corresponding real frequency and reduces the wavenumber at maximum growth (Winske, 1988). The figure is redrawn from Figures 2, 3 in Winske (1989), with permission from The American Institute of Physics.

rate γ_m and corresponding wavenumber k_m at maximum growth are given by

$$\gamma_m \approx \omega_r \approx k_m V_n \approx \omega_{LH}, \quad (11)$$

where $\omega_{LH} = \omega_{pi}/(1 + \omega_{pe}^2/\omega_{ce}^2)^{1/2} \approx \omega_{ci}(m_i/m_e)^{1/2}$. When $V_g \gg V_n$, the instability reduces to the LLR solution of Hassam and Huba (1987).

Figure 5 compares solutions a simplified linear dispersion relation (electrostatic, cold electrons, cold ions) showing the real and imaginary part of the frequency vs. wavenumber for two cases: blue curves ($V_n/v_A = 1$, $V_g = 0$) and red curves ($V_n/v_A = 1$, $V_g = 3$) (Figures 2, 3 in Winske, 1989). Here the ion beta is assumed to be $\beta_i = 0.2 = v_i^2/v_A^2$, with v_i the ion thermal speed and v_A the Alfvén speed. With $V_g = 0$ the wave properties are those expected of the lower hybrid drift instability, i.e., high frequency, fast growing, short wavelength unstable modes. Whereas, with the inclusion of $V_g = 3 v_A$, the instability grows even faster, with maximum growth occurring at longer wavelength. It should be noted that from Cavity Formation and Properties, the deceleration of the expanding plasma increases as the plasma expands and reaches its maximum as the plasma stops. In addition, the density gradient at the leading edge can also change in time. Thus, when doing linear theory for specific experimental conditions, such as the AMPTE releases, one typically uses average values for the plasma parameters, rather than trying to determine more realistic local values during the plasma expansion. In addition to appearing as the observed flute modes, the unstable waves also heat both electrons and ions, thus leading to the thermalized state of the expanded plasma

when it comes to rest as the magnetic cavity reaches its maximum size.

Huba et al. (1990) generalized the kinetic local linear theory to include additional effects, namely ion-neutral collisions and related Pedersen drifts, generalized magnetization effects of the ions through the use of Gordeyev integrals, and electromagnetic effects on the electrons. They also employed proper electron and ion distribution functions in setting up the equilibrium. Numerical solutions, in which the role of the gravitational drift were emphasized, were presented as well as analytic solutions of the dispersion equation in various limits. Parameters for the AMPTE and CRRES barium releases in space as well as for the NRL laser experiment were carefully compiled and used in the linear theory calculations that were then compared with measurements. Good qualitative agreement was found in all cases, but generally, the linear theory gives wavelengths of the fastest growing modes that are too short (often by an order of magnitude) compared with observations.

The linear theory of these fast-growing instabilities augmented by the gravitational drift was also extended to include non-local effects to provide a more realistic picture of the interface between the expanding plasma cloud and the magnetic field. Huba et al. (1989) modified the Hall-MHD model and Gladd and Brecht (1991) extended the kinetic model. In both cases eigenfunctions were obtained that characterize the radial structure of the electric field at the interface. Peak growth rates obtained from the non-local analysis were also consistent with those obtained from local theory, again emphasizing that linear theory tends to overestimate the growth rates of the most unstable waves.

Simulations

A number of 2-D simulations of the development of striations on expanding plasma clouds were also carried out, using a variety of plasma models to both show the validity of the linear theory calculations and also to examine the non-linear behavior. (We defer showing examples of the instability development in the simulations to later in the section when recent 3-D simulations will be discussed and compared with the earlier 2-D calculations). Huba et al. (1987) conducted MHD and Hall-MHD simulations in a slab geometry (denser fluid supposed by lighter fluid and stronger magnetic field) with an imposed “gravity.” In the MHD simulations, the usual “spike and bubble” structure appears. In the Hall-MHD simulations, the character of the modes is much different, i.e., shorter wavelength, faster-growing, more finger-like structures appear (also see Winske, 1996). Similar MHD and Hall-MHD simulations (Huba et al., 1992, 1993) were also done later for the G-4 and G-10 barium releases associated with the CRRES mission. No striations appeared in the simulations using the MHD equations, whereas flute modes rapidly formed when the Hall-MHD term was added to the calculations. Also in these simulations the flutes evolved to longer wavelengths in time while the tips of the flutes seemed to continue to propagate outwards even when most of the barium ions had stopped.

Similarly, full-particle 2-D simulations of expanding plasma clouds were carried out by Winske (1988, 1989). Consistent with linear theory, very short wavelength modes appeared and grew on lower-hybrid time scales (e.g., Figure 11 in Winske, 1988). Analysis of the wave amplitudes indicated that growth during the linear stage was consistent with linear theory and wave generation occurs at very short wavelength. However, it should be pointed out that such waves could have been seeded on the expanding cloud at early times by the computational grid. In addition, the saturation of the waves was much larger than expected from usual arguments for the lower hybrid drift instability (Liewer and Davidson, 1977). In this case the continuing expansion of the plasma cloud keeps the radial electric field, which gives rise to the azimuthal current, large so that the system continues to be driven to much higher wave levels that persist until the plasma stops expanding. Some coalescence to longer wavelengths appeared after the plasma stopped and the instability saturated. In various simulations, as ρ_d/R_B was increased, the size of the flutes became larger and the final cavity size was smaller (Winske, 1989). These results were consistent with the experimental results of Dimonte and Wiley (1991) and Zakharov et al. (2006). A number of other full particle simulations were also carried out by various groups, using different initial conditions and plasma parameters (see review by Akimoto et al., 1988). Such calculations demonstrate the robustness of the instability generation mechanism.

Hybrid simulations, in which the electrons are treated as a massless fluid, the ions are treated kinetically using particle-in-cell techniques and electromagnetic fields are considered in the low-frequency approximation (e.g., Winske and Gary, 2007), were also been carried out for expanding plasma clouds in both 2-D and 3-D geometries (Brecht and Thomas, 1988; Brecht and Gladd, 1992). Because of the lack of electron inertia, linear theory in this case indicates that the linear growth rate of surface flute

modes increases with wavenumber out to the resolution of the calculation. In this case it is necessary to initialize the simulations with a prescribed short wavelength perturbation. As the plasma expands, this mode grows at a rate consistent with linear theory and continues to dominate to saturation. Later, there may be coalescence to longer wavelengths. Such simulations provide an interesting compromise between Hall-MHD and full particle simulations. Hybrid simulations with a finite electron mass have also been carried out in a 2-D slab geometry (in which case the deceleration is slower, being constant in time and space) that show mode coalescence at late times that increases in simulations where the ion to electron mass ratio is increased (Sgro et al., 1989).

Summary

By the early 90's, the experimental programs were essentially finished and most of the modeling had also been completed. The consensus among those who had worked on this problem could be summarized in a schematic figure, **Figure 6** (from Figure 17 in Huba et al., 1990). The figure shows the (random) development of short wavelength modes as the plasma cloud compresses into a thin, expanding shell. By the time of saturation, the instability had evolved to longer wavelength modes. At even later time, these evolved modes continue to grow and expand outward, even as the inner edge of the cavity begins to collapse. This picture was consistent with experiments and simulations at that time, but evidently did not provide a detailed explanation. Zakharov et al. (1986) proposed a mechanism to explain why the cavity seemed smaller and the flutes larger when ρ_d/R_B was larger, based on the competition between expansion and diffusion. Qualitatively this model is consistent with experiments and simulations, which

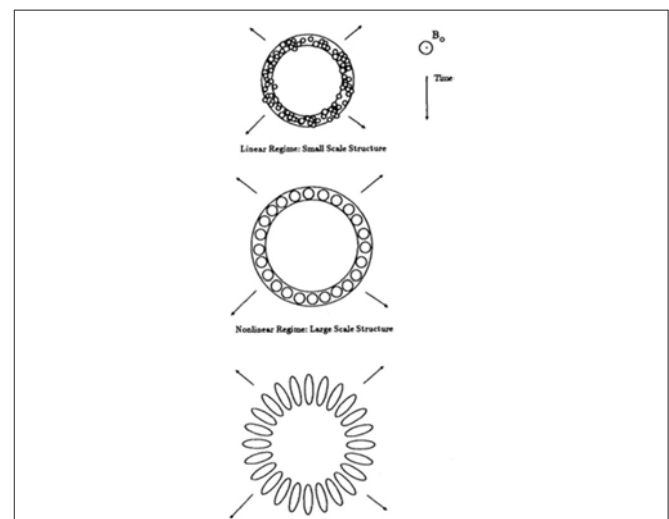


FIGURE 6 | Pictorial representation of the understanding of striation formation and development from Figure 17 in Huba et al. (1990), showing progression from small random waves generated in the expanding shell at early times, coalescence to larger wavelengths and then expansion of the longer wavelengths modes beyond the surface of the shell at later times. The figure is reproduced with permission from The American Institute of Physics.

indicate that the instability, and the resulting diffusion, are larger when ρ_d/R_B is large (Winske, 1989).

New Work

In the last few years new experiments and simulations have shed further light on the processes of flute mode generation and non-linear development. For example, experiments at the University of California Los Angeles (UCLA) by Collette and Gekelman (2011) used a much smaller laser, ~ 1.2 J, but one which could be fired in sync with the background plasma that is produced at 1 Hz in the Large Plasma Device (LAPD) (Gekelman et al., 1991). In addition, small computer-controlled probes that measure the components of the magnetic field could be moved automatically every few shots so that the dynamics of the cavity formation could be carefully mapped out from the B field components and the current J (computed from $\nabla \times B$) in space and time. For example, the current structure that produces the diamagnetic cavity inferred from such measurements is shown in the right panel of **Figure 7** (Figure 6 in Collette and Gekelman, 2011). Even though the cavity sizes in these experiments were small ($\sim$ few cm), the movable probe measurements can be converted into movies and fast photographs can be used to measure the growth and development of flutes, as shown in the left stack of panels in **Figure 7** (Collette and Gekelman, 2011, Figure 2). These experiments demonstrate that while they are in a somewhat different physical regime, in terms of dimensionless parameters they are consistent with earlier space and laboratory experiments. The very high-resolution results show the development and evolution of very short wavelength modes that coalesce into larger structures. The presence of collisional effects in these experiments could also be quantified and indicate that the collapse of the cavity occurs much faster than expected from collisional diffusion.

In addition to new experiments, new simulations of diamagnetic cavities have also been recently carried out. For example, Huba has performed new 3D ideal and Hall MHD simulations of an expanding ion cloud. For these simulations the ions are deposited using

$$\frac{\partial n_i}{\partial t} = n_0 \sigma_i [1 - \exp(-\sigma_i t)] \exp(-\sigma_i t) \exp(-\{(r - V_0 t)/\Delta r_0\}^2), \quad (12)$$

where $n_0 = 5 \times 10^4 \text{ cm}^{-3}$ is the initial neutral density, $\sigma_i = 0.14 \text{ s}^{-1}$ is the ionization rate, $V_0 = 1 \text{ km/s}$ is the expansion velocity, and $\Delta r_0 = 0.5 \text{ km}$. The simulation is initialized at $t = 2 \text{ s}$ with an initial spherical shell of expanding ions at a radius $r = 2 \text{ km}$. Additionally, there is a 1D background magnetic field $B_z = 10^{-4} \text{ G}$. These parameters are chosen for numerical expediency and do not represent a specific magnetospheric barium release. The grid used is $100 \times 100 \times 100$ and is non-uniform; the extent of the grid perpendicular to B is $\pm 55 \lambda_i$ and along B is $\pm 140 \lambda_i$, where λ_i is the ion inertial length.

The results are shown in **Figures 8, 9**. We show an isodensity contour $n_i = 700 \text{ cm}^{-3}$ at times $t = 15 \text{ s}$ and 32 s for the ideal MHD case (left panels) and Hall MHD case (right panels). The z -axis is aligned with the background magnetic field. In **Figure 8** the view of the contours is along the z -axis, i.e., along

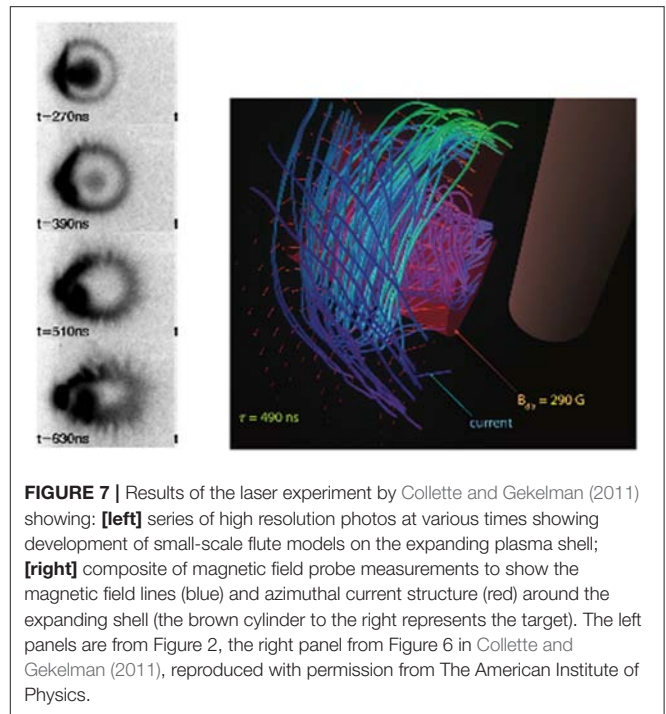


FIGURE 7 | Results of the laser experiment by Collette and Gekelman (2011) showing: **[left]** series of high resolution photos at various times showing development of small-scale flute models on the expanding plasma shell; **[right]** composite of magnetic field probe measurements to show the magnetic field lines (blue) and azimuthal current structure (red) around the expanding shell (the brown cylinder to the right represents the target). The left panels are from Figure 2, the right panel from Figure 6 in Collette and Gekelman (2011), reproduced with permission from The American Institute of Physics.

the magnetic field. At $t = 15 \text{ s}$ both the ideal and Hall MHD cases are essentially identical. However, at $t = 32 \text{ s}$ the ideal MHD case is basically unstructured, while the Hall MHD case has become unstable and large-scale density irregularities have developed. In **Figure 9** the view of the contours is perpendicular to the background magnetic field. Again, at $t = 15 \text{ s}$ both the ideal and Hall MHD cases are essentially identical, but the contours are more extended along the magnetic field because the ions can freely expand along this axis. At $t = 32 \text{ s}$ the ideal MHD case is not structured and the extension along the magnetic field is more pronounced. The Hall MHD case shows the development of density irregularities that are filamented along the magnetic field, as expected.

The earlier full particle simulations (Winske, 1988, 1989) of short wavelength flute modes developing on the surface of an expanding debris plasma cloud were carried out in two spatial dimensions perpendicular to the background magnetic field. The observed wavelengths were in qualitative agreement with the linear theory of the generalized lower hybrid drift instability. The initial wavelengths were very short, perhaps limited by the grid resolution, but they developed into somewhat longer wavelength structures in time. While it was expected that there would be some changes if the simulations were run in 3-D rather than 2-D, i.e., the deceleration of the cloud is stronger in 3-D and develops later in the expansion as shown in **Figure 5**, it was not possible to carry out 3-D simulations of even modest size three decades ago. With the development of modern 3-D particle codes, e.g., VPIC (Bowers et al., 2008), and the readily available computing resources, 3-D simulations have now been done and compared with 2-D, as shown in **Figure 10**. The parameters of the simulation are very similar to those in Winske (1988); the debris

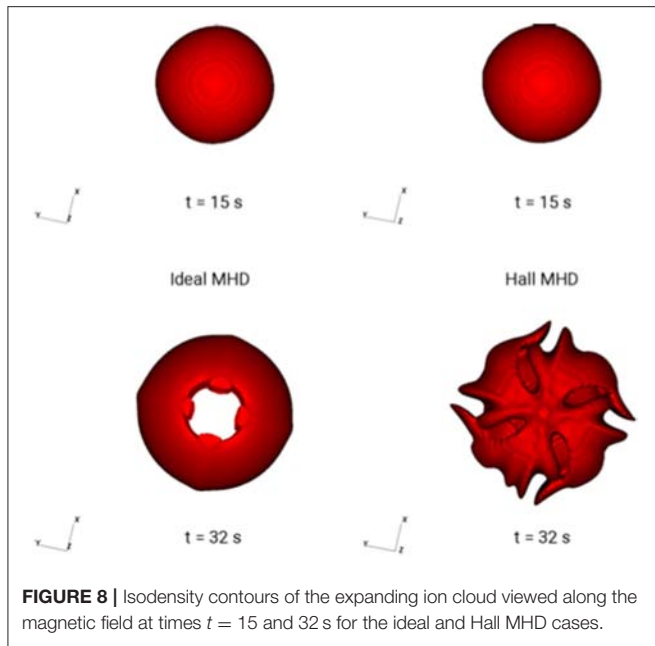


FIGURE 8 | Isodensity contours of the expanding ion cloud viewed along the magnetic field at times $t = 15$ and 32 s for the ideal and Hall MHD cases.

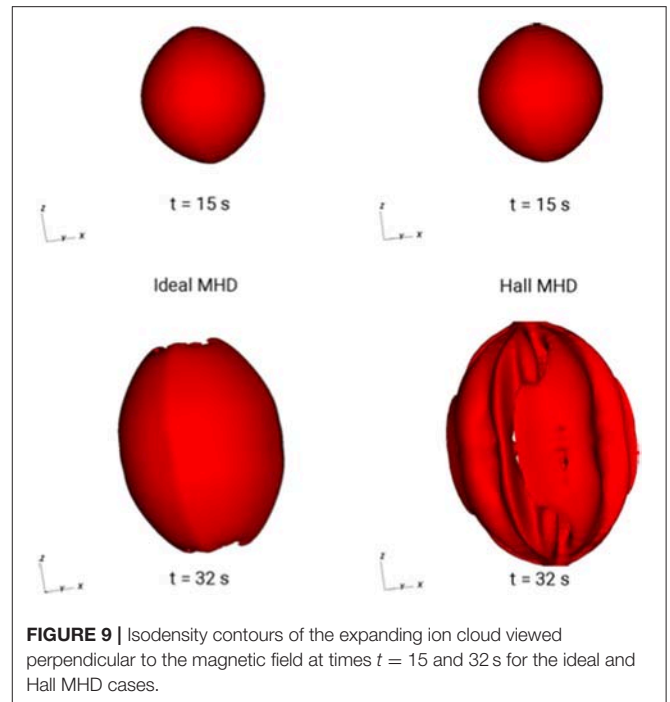


FIGURE 9 | Isodensity contours of the expanding ion cloud viewed perpendicular to the magnetic field at times $t = 15$ and 32 s for the ideal and Hall MHD cases.

ions have $m_d/m_e = 100$, initial cloud density $n_d/n_o = 25$, initial radius $= 2 c/\omega_{pe}$, expansion velocity $v_d = v_A$, $\omega_{pe}/\omega_{ce} = 5$. The system size is $25 \times 25 c/\omega_{pe}$ in 2-D, using 5 million simulation particles to represent the debris ions and electrons and 25 million particles each to present the background ions and electrons. Similarly, the system size is $25 \times 25 \times 25 c/\omega_{pe}$ in 3-D, using ~ 1 billion simulation particles to represent the debris ions and electrons and 8.3 billion particles each to present the background ions and electrons. **Figure 10** shows a summary of the two simulations. The left side of the figure shows color contours of the debris electron density in 2-D (left panels) and 3-D (right panels) at various times. At the early time ($\omega_{ce}t = 25$, corresponding to $\omega_{LH}t \approx 2.5$), the instability in 2-D has begun to grow. By $\omega_{ce}t = 50$, the instability is well-developed in 2-D and just starting to grow in 3-D. By $\omega_{ce}t = 75$ the instability is fully developed in 2-D with well-defined, but low-density flute modes extending beyond the radius of the cavity. In 3-D the instability is not yet as well-developed, but flute modes growing at slightly longer wavelength are apparent. The panel on the right of **Figure 10** shows the 3-D image of the debris electron density at $\omega_{ce}t = 75$ that more clearly reveals the field-aligned striations on the surface and how they would appear to an external viewer. The black “fuzz” represents a low-density halo of electrons that have been able to propagate away from the outer edge of the cavity. Overall, these comparative calculations are consistent with expectations: linear theory would suggest that the instability should develop somewhat slower in 3-D at slightly longer wavelengths due to the difference of the deceleration, $g(t)$. And from the 3-D perspective the striations on the surface are reminiscent of the AMPTE image (**Figure 1**). Evidently, more quantitative comparisons from these simulations need to be carried out to examine what are the levels of the fluctuating electric fields and how the linear modes coalesce in time.

EXTENSIONS

Finally, we consider more recent work related to magnetic cavity formation. In this case we consider expansion of a dense target plasma across a magnetized background plasma, but with the expansion speed greater than the Alfvén speed. As we show, the physics becomes different as the interaction is dominated by the dynamics of the background plasma as the cavity forms and this interaction determine the cavity size.

Recall that in the initial discussion of the cavity size, the energy equation, Equation (1), was used, assuming that the background density was so low that it could be ignored. Here the opposite limit is considered, where the background density is large and its contribution dominates that of the magnetic field. Instead of Equation (1), conservation of momentum is used to determine the distance over which the expanding target mass M overruns an equivalent mass of background plasma (ion mass m_i , ion density n_o), the so-call equal mass radius R_m :

$$M = m_i n_o \frac{4\pi R_m^3}{3}. \quad (13)$$

Comparing the equal mass radius with the magnetic confinement radius, R_B , determined from Equation (1), we find

$$\frac{R_m^3}{R_B^3} = \frac{v_A^2}{V_d^2}. \quad (14)$$

where v_A is the Alfvén speed based on the background plasma density and the magnetic field. It is evident from Equation (14) that the expanding plasma will stop (and hence determine the size of the magnetic cavity) at $R = R_B$ if the background plasma

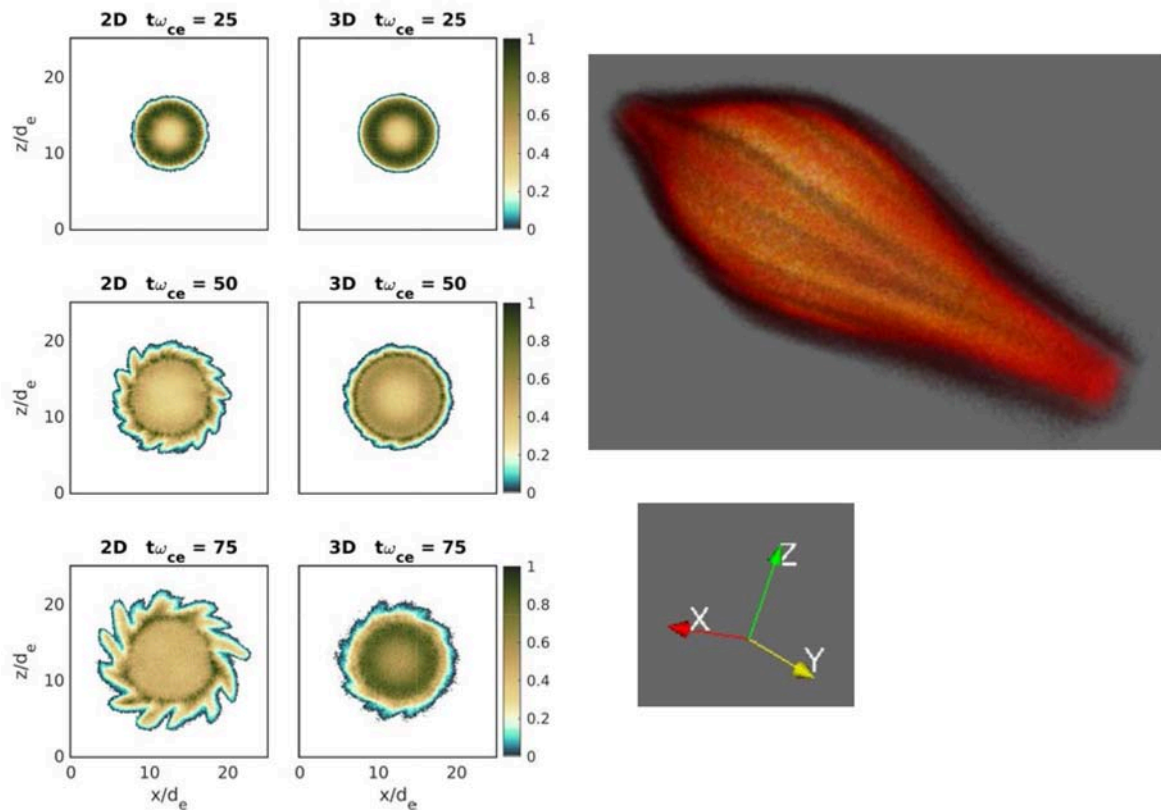


FIGURE 10 | Results of 2-D and 3-D VPIC particle simulations: left panels show 2-D debris electron density in the plane perpendicular to the magnetic field from (left) 2-D run and (right) 3-D run at three times (top row – $\omega_{ce}t = 25$, when the instability just begins to develop in the 2-D run; middle row – $\omega_{ce}t = 50$, when instability is fully developed in 2-D but just starting in 3-D; lower row – $\omega_{ce}t = 75$, when instability is fully developed). The right panel shows a 3-D picture of the debris electron cloud surface at this later time, $\omega_{ce}t = 75$, showing the cloud expansion along the magnetic field (y-direction) and how the surface perturbations would appear to an external viewer.

density is low so that the expansion speed is less than the Alfvén speed. Similarly, if the background density is larger and thus the Alfvén speed is smaller, the expanding plasma stops, and the cavity radius is determined, by the equal mass radius, R_m . As we discuss later, in this case the expanding debris ions transfer their momentum to the background ions through a process known as Larmor coupling. Another important aspect of cavities produced by super-Alfvén Mach number expansions ($V_d > v_A$) is that there are no flute modes generated on the cavity surface, as we discuss later.

The AMPTE mission conducted heavy ion releases in solar wind in front of and just behind the bow shock. In this case the solar wind flow is significantly larger than the Alfvén speed so that in the solar wind frame the photoionized barium and lithium ions are super-Alfvénic. Observations of these events (Lühr et al., 1986; Valenzuela et al., 1986) do show cavity formation, but with very small cavities resulting from the lithium releases. But more significant observations involved the subsequent motion of the plasma cloud in the solar wind after the cavity collapses, for example, the barium release of Dec. 27, 1984. There was some dispute about the initial motion of the cloud transverse to the solar wind flow, which Papadopoulos et al. (1987) explained in terms of the motional electric field in the solar wind. There was

also considerable interest about the subsequent behavior of the lithium ions that were picked up by the solar wind and formed an artificial comet (Valenzuela et al., 1986).

In recent years laboratory experiments have also been done using high-power lasers to generate high-energy target ions that expand at super-Alfvénic speeds across a background magnetic field, both at UCLA on LAPD using the > 100 J Raptor laser (Niemann et al., 2012) and at the Laser Institute in Novosibirsk. Experiments and simulations (Clark et al., 2013, 2014; Niemann et al., 2013, 2014; Schaeffer et al., 2014, 2015; Shaikhislamov et al., 2015) exhibited cavity formation and development. More importantly, they showed that a collisionless shock wave is generated that separates from the edge of the cavity and propagates outward. In the super-Alfvénic case, the pickup of the background ions allows the momentum of the outward streaming target ions to be transferred to the background ions. Equation (13) indicates that the expanding target ions transfer their momentum and energy to the background ions to conserve momentum, but does not indicate how this occurs. Because these are collisionless plasmas, the transfer cannot occur through collisional interactions, but instead involves the electromagnetic fields. In particular, the coupling occurs through the transverse electric field (E_ϕ), which is the same mechanism

that allows newborn ions in the solar wind to be “picked up,” i.e., accelerated up to the solar wind speed to form an artificial comet (Papadopoulos et al., 1987). Because of the presence of the electric field, the background ions are initially accelerated in the azimuthal direction, and then by gyromotion in the magnetic field acquire a large positive radial velocity, as the radially expanding debris ions also start to bend in the magnetic field and decrease their radial velocity (eventually with their radial velocity going to zero). This mechanism, known as Larmor coupling, has been investigated by Bashurin et al. (1983), Hewett et al. (2011) and more recently by Bondarenko et al. (2017b). It has also been recently demonstrated in the laboratory using ion spectroscopy (Bondarenko et al., 2017a) as well as with ion current probes (Prokopov et al., 2016). It is to be noted that because the background ions continue to carry momentum and energy outward, at the cavity boundary the magnetic disturbance does not stop but continues propagating outward as a collisionless shock wave.

The formation of the magnetic cavity in the UCLA experiments is measured using magnetic field probes. Five probes at different radial positions are mounted on a rod, which is moved each time the laser is fired (once every 45 min) in order to map out in space and time the magnetic field intensity. The top panel

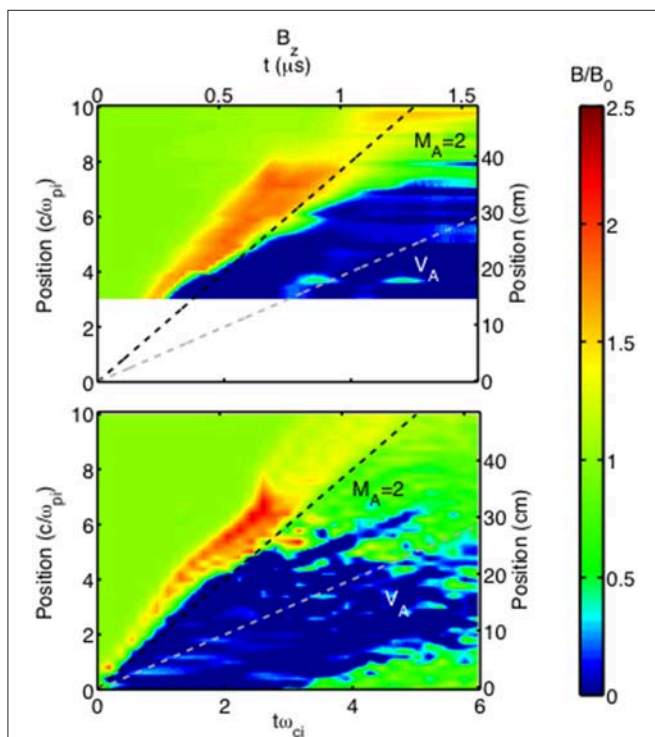


FIGURE 11 | Comparison between experimental data from the UCLA laser experiment (top) and a 2-D hybrid simulation (bottom). The plots show the magnetic field magnitude, B/B_0 , as a function of time and space. The formation and evolution of the magnetic cavity (blue region) and the development of a collisionless shock at the edge of the cavity (red) are clearly visible in each plot and show good agreement between simulation and experiment. The plot is from Figure 5 of Clark et al. (2014) and is reproduced with permission from The American Physical Society.

in **Figure 11** shows one such compiled figure (Figure 5, Clark et al., 2014). Values of the magnetic field B/B_0 at various probe positions are plotted as a function of time (with positions and time given both in actual and normalized units). (Near the target there are no measurements). Two features stand out. The dark blue region, corresponding to $B \sim 0$, is the developing magnetic cavity, which eventually expands out to $r \sim 35$ cm, corresponding to $\sim 7 c/\omega_{pi}$. The red region is the compressed magnetic field that forms early in time and propagates away from the cavity at Alfvén Mach number $M_A \sim 2$.

Simulation is an important tool in these experiments, since at present there are no diagnostics to measure the plasma properties of either the debris ions or the background ions. Two-dimensional hybrid simulations (Clark et al., 2013, 2014; Niemann et al., 2013, 2014; Schaeffer et al., 2014, 2015) have been carried out for the conditions in the experiments and generally are in good agreement. The bottom panel of **Figure 11** shows the comparison of the simulations (Figure 5, Clark et al., 2014) for the experimental conditions in the top panel, plotting the magnetic field magnitude vs. space and time in the same format on the same scales. The simulation reproduces the size of the cavity and the speed of the outgoing magnetic wave. The measured magnetic field change across the wave front can be shown to be consistent with the jump conditions that are satisfied for a collisionless shock wave. The 2-D simulations of the experiment also show that as the cavity develops, flute modes do not appear on the surface and some of the target ions have diffused beyond the cavity boundary (Clark et al., 2013, 2014). A number of experiments of this type have been carried out with different initial conditions that help define the criteria for which a collisionless shock can be generated (Schaeffer et al., 2017a).

It is also possible to redirect the laser beam so that it enters the plasma chamber at a small angle with respect to the axial magnetic field. By repositioning the target closer to one end of the plasma column, one can then study the dynamics of energetic debris ions streaming along the magnetic field. As before, the expansion of the energetic debris ions from the target produces

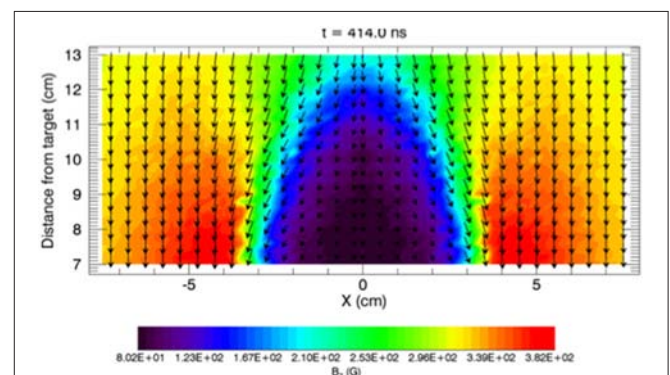


FIGURE 12 | Magnetic probe measurements of the magnetic field direction and magnitude (arrows) produced by a low Mach number ($M_A \sim 1$) expansion of target ions with laser beam oriented along the background magnetic field. The magnetic cavity (purple region) forms just in front of the target (Heuer et al., 2018, private communication).

a magnetic cavity, which can be measured by magnetic probes, as shown in **Figure 12** (Heuer et al., 2018, private communication). In this case a small (~ 20 J) laser that fires at the same rate that the background plasma is generated is used to make the cavity, allowing very detailed probe measurements in 3-D. The cavity has a small transverse scale (radius ~ 3 cm) but is greatly expanded along the field (~ 12 cm). The interaction of the debris ions with this cavity tends to scatter a majority of the ions outward toward the walls of the chamber where they are lost. But a fraction of the debris ions do continue to propagate along the magnetic field at several times the Alfvén speed and generate electromagnetic waves (Heuer et al., 2018). Magnetic probes at various positions downstream from the target measure the magnetic wave components along and transverse to the ambient field. Hybrid simulations and theory are then used to compare the computed wave properties with the measurements (Weidl et al., 2016); good agreement is obtained.

Lastly, it should be pointed out that other types of magnetic cavities occur in space near the bow shock. These structures were originally called hot diamagnetic cavities (Thomsen et al., 1988) or hot flow anomalies (Schwartz et al., 1988) and were thought to be relatively small-scale entities just upstream of the shock. These structures result from interactions of discontinuities or steepened waves in the solar wind with the bow shock. The disruption of the shock allows downstream ions to flow into the solar wind, somewhat similar to a chemical ion release upstream of the bow shock, creating a cavity-like structure. These are not genuine “diamagnetic cavities” in the sense that the expanding, hotter plasma density from the magnetosheath is not much greater than the background (solar wind) density and thus the magnetic field is not totally excluded inside the cavity. The size of these cavities can vary widely. More recent observations (with simultaneous multi-spacecraft measurements) and 2-D hybrid simulations indicate they can extend over significant portion of the bow shock (Omidi et al., 2013).

CONCLUSIONS

The formation and properties of diamagnetic cavities were very active areas of space and laboratory plasma research several decades ago with both ion release experiments in space and laser-produced ion expansion experiments in the laboratory. Both types of experiments demonstrated magnetic field exclusion by the expanding release/target generated plasma and subsequent cavity formation. They also showed formation of field-aligned striations on the cavity surface, with wavelengths of the flute modes less than the ion gyroradius and much smaller than the cavity size. Linear theory indicated that a generalized lower-hybrid instability driven by both the density gradient at the edge of the cavity and more importantly, an additional drift due to the deceleration of the cavity surface could be the source of such unstable flute modes. Hall-MHD and particle simulations verified the instability mechanism, although there were issues concerning the observed wavelengths being much longer than the most unstable modes predicted by linear theory. By the early 90's, when the experimental campaigns were winding down, this was the state of our understanding (i.e., **Figure 6**).

Since then, there have been some additional laboratory experiments of sub-Alfvénic plasma expansions with improved diagnostics. But most importantly, very recent 3-D Hall-MHD and full particle simulations more clearly reveal that the initial development of the instability does occur at very short wavelength—perhaps even determined by grid resolution in the calculations—but the instability rapidly evolves to somewhat longer wavelengths, more consistent with the experimental measurements. Even though the effect of the deceleration of the plasma is a stronger effect in 3-D compared to 2-D, and the (linear) instability behavior is consistent with this difference, overall there are not new effects that arise in three dimensions. This was not totally unexpected since for the phenomena that we are interested in, namely cavities and surface striations, we have already demonstrated in this review that the space experiments, which are three-dimensional, are similar in many ways to the laser experiments that are essentially two-dimensional.

Most recent experimental and computational research has been devoted to debris ion expansions at greater than the Alfvén speed as well as expansions along, rather than across the magnetic field. In these cases, cavities are still produced, but instead of flute modes at the surface, collisionless shock waves propagate outward from the cavity when the expansion is primarily perpendicular to the ambient magnetic field and various types of ion beam-driven electromagnetic waves are produced in parallel expansions. While space observations reveal many properties of shocks and related particle acceleration that can be studied in the laboratory. Howes (2018), the laboratory experiments are unique for clarifying the time-dependent processes of how cavities form and shocks are generated. In the future it may be possible to conduct such experiments in larger facilities with imposed inhomogeneous background plasmas and magnetic fields. For example, experiments at the Omega EP laser facility have already demonstrated that high-Alfvén-Mach number shocks can be generated (Schaeffer et al., 2017b). Interesting new results from such experiments concerning cavity sizes and shapes, as well as wave generation and non-linear development, verified by appropriate 2-D and 3-D Hall-MHD, full particle and hybrid simulations, could eventually lead to new interest for active experiments in space. One promising candidate is a proposed small barium release via a rocket launch to generate lower hybrid turbulence and whistler waves (Ganguli et al., 2015).

AUTHOR CONTRIBUTIONS

DW wrote much of the text and put the manuscript together. He was responsible for putting together 7 of the figures and related text. JH provided results of new calculations, resulting in 2 figures plus related text. CN provided 2 figures and related text. AL provided results of new calculations, resulting in 1 figure plus related text.

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Review of Controlled Excitation of Non-linear Wave-Particle Interactions in the Magnetosphere

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Controlled experiments involving injection of 0.5 Hz–8 kHz electromagnetic waves into the Earth's magnetosphere have played an important role in discovering and elucidating wave-particle interactions in near-Earth space. Due to the significant engineering challenges of efficiently radiating in the ELF/VLF: 300 Hz–30 kHz band, few experiments have been able to provide sustained transmissions of sufficient power to excite observable effects for scientific studies. Two noteworthy facilities that were successful in generating a large database of pioneering and repeatable observations were the Siple Station Transmitter in Antarctica and the High Frequency Active Auroral Research Program (HAARP) facility in Alaska. Both facilities were able to excite Doppler shifted cyclotron resonance interactions leading to linear and non-linear wave amplification, triggering of free running emissions, and pitch angle scattering of energetic electrons. Amplified and triggered waves were primarily observed on the ground in the geomagnetic conjugate region after traversal of the magnetosphere along geomagnetic field aligned propagation paths or in the vicinity of the transmitter following two traversals of the magnetosphere. In several cases, spacecraft observations of the amplified and triggered signals were also made. The observations show the amplifying wave particle interaction to be dynamically sensitive to specific frequency and also specific frequency-time format of the transmitted wave. Transmission of multiple coherent waves closely spaced in frequency showed that the wave particle interaction requires a minimum level of coherency to enter the non-linear regime. Theory and numerical simulations point to cyclotron resonance with counter streaming particles in the 10–100 keV range as the dominant process. A key feature of the non-linear interaction is the phase-trapping of resonant particles by the wave that is believed to drive non-linear wave amplification and the triggering of free-running emissions. Observations and modeling of controlled wave injections have important implications for naturally occurring whistler mode emissions of hiss and chorus and the broader phenomena of radiation belt dynamics. A review of observational, theoretical, and numerical results is presented and suggestions for future studies are made.

Keywords: whistler anisotropy instability, triggered emissions, whistler mode chorus waves, active experiments, HAARP facility, radiation belts, space weather

1. HISTORY AND SIGNIFICANCE OF WHISTLER MODE OBSERVATIONS AND ACTIVE EXPERIMENTS

Appreciation of the role of whistler mode waves in the near-Earth space environment predates the space age and began with the landmark publication by Storey (1953). Storey (1953) identified the plasma nature of the space around of the Earth (out to several Earth radii of altitude) as responsible for the phenomena of “whistling” atmospherics which were first observed on communications hardware during World War I. He described how lightning induced impulsive radiation in the ELF/VLF band (3 Hz–30 kHz) couples through the ionosphere, into the magnetosphere and experiences frequency dispersion due to propagation along the geomagnetic field line in a right hand circularly polarized mode below the electron cyclotron frequency and plasma frequency. This mode has since been called the whistler mode. The work of Storey (1953) was foundational in magnetospheric physics in that it not only established the magnetosphere as filled with significant densities of cold plasma, but also was the first to describe natural whistler wave emissions, known as hiss and chorus, subsequently identified to result from hot plasma instabilities. Today whistler mode waves of both terrestrial and magnetospheric origin are seen as key drivers in near-Earth space energy dynamics (Reeves et al., 2003; Bortnik and Thorne, 2007; Thorne, 2010). Despite several decades of research, whistler-mode wave particle interactions continue to be the subject of active investigations since the near-Earth space environment and its energy dynamics are of increasing economic and strategic importance. Recently there has been renewed interest in non-linear whistler mode phenomena and a consensus that the non-linear regime of wave-particle interactions needs to be quantified to achieve accurate prediction capabilities in global flux and energy models. Active whistler mode injection experiments, which are the topic of this review, have been the drivers of non-linear phenomena investigation and can play an important role in future efforts.

Almost a decade after Storey’s results (Storey, 1953) were published, it was discovered that it is possible to actively trigger whistler mode emissions in the magnetosphere with controlled transmissions from VLF communication transmitters (Helliwell et al., 1964). A VLF receiver on board a ship USNS *Eltanin* in the magnetic conjugate region of the U.S. Navy NAA transmitter in Cutler, Maine observed amplification and triggering of new frequencies from the Morse code 14.7 kHz transmissions. The observed records revealed that the emissions were triggered almost exclusively by the 150 ms Morse dashes and only rarely by the 50 ms Morse dots (Helliwell et al., 1964; Helliwell, 1965, p. 297–298). This remarkable phenomenon was dubbed the “dot-dash anomaly” and sparked interest in dedicated transmissions at variable frequency for controlled experiments of magnetospheric wave particle interactions.

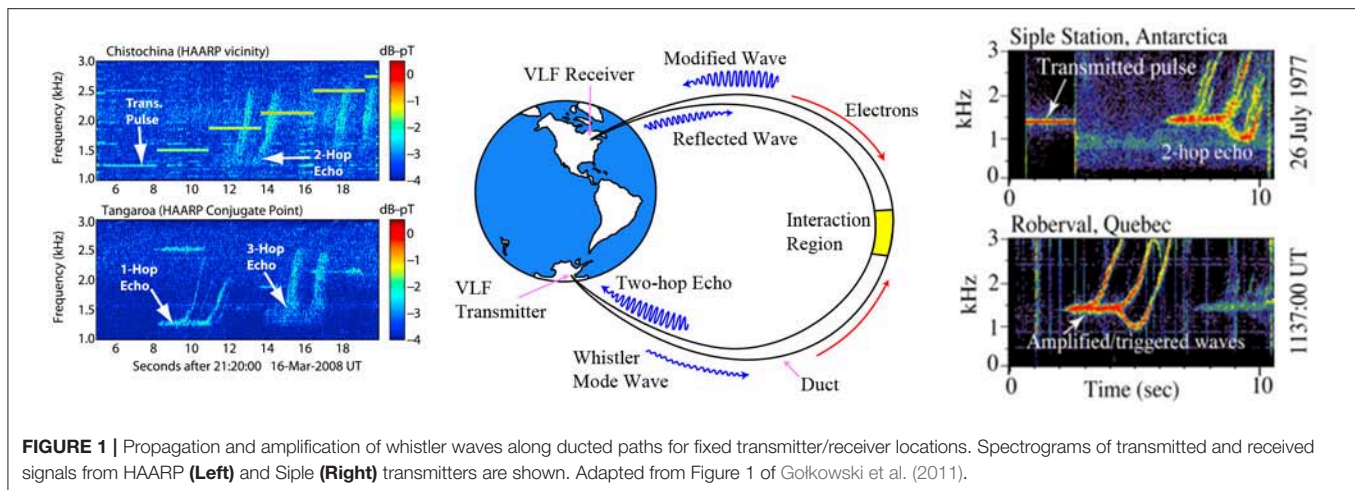
As illustrated in **Figure 1**, an hemisphere to hemisphere wave injection experiment turns the inner magnetosphere into a plasma chamber in which controlled whistler mode sources can be used as diagnostics of the condition of the plasma and for excitation of instabilities. In such experiments transmitted

signals which have made one traverse through the magnetosphere and are observed in the magnetic conjugate region are known as “one hop echoes” and signals that have made two traverses and returned to the transmitter region are known as “two hop echoes.” Antarctica was initially seen as an optimal location for a VLF wave injection experiment where transmissions into the magnetosphere along geomagnetic field lines would be possible. The advantages of Antarctica included the lack of major sources of electromagnetic noise including man-made interference and thunderstorm activity, the established observations of natural VLF emissions, and the accessibility of the conjugate locations on landmasses in the northern hemisphere. Furthermore, the presence of thick ice sheets allowed for significant elevation of an horizontal antenna above the conducting surface of the Earth. An initial attempt of a transmitter near Byrd Station (80.02° S, 119.53° W, $L \sim 7.2$) known at the Byrd Longwire was operated from 1966 to 1969 but yielded mixed results (Helliwell and Katsufakis, 1974; Gibby, 2008). One reason that the Byrd Longwire was not able to excite signals that could be observed at the conjugate point was that it was located at a high L shell where geomagnetic field lines can be open and hemisphere to hemisphere ducting is unfavorable. Whistlers in Antarctica were typically observed to have propagated along paths near $L \sim 4$.

1.1. Siple Station

In 1969 an effort was undertaken to find an ideal site in Antarctica for a VLF wave injection experiment and after an exhaustive search, a site was selected at 79.93° S, 84.25° W, 2381 km east of McMurdo Station. At $L \sim 4.3$, the site, named Siple Station in honor of American Antarctic pioneer Paul Siple, offered access to high magnetic latitudes, the plasmapause, and natural VLF emissions. The magnetic conjugate point was located near the city of Roberval in Quebec, Canada, making establishment of a conjugate monitoring station straightforward. The initial installation, completed in 1973 used the 80 kW VLF transmitter from Byrd Station and a single 21.2 km horizontal antenna giving a resonant frequency of approximately 5 kHz. The installation received significant upgrades over the years with a 150 kW transmitter installed in 1979, the antenna lengthened to 42 km in 1983 and the addition of a second 42 km dipole in 1986. The 42 km crossed dipoles of the final installation were resonant at 2.5 kHz and could directly excite a right hand polarization that could propagate in the ionosphere and magnetosphere (Helliwell, 1988).

Analysis by Raghuram et al. (1974) showed that the antenna efficiency at Siple Station was on the order of 2–3%. The 2 km thick ice sheet was key in elevating the antenna above the conducting ground and mitigating detrimental image currents (Helliwell, 1988). This controlled science dedicated injection of several kW of power in the few kHz band continues to be unmatched to this day. The Siple experiment was very successful in producing observations of non-linear growth and triggering of whistler mode waves in the conjugate region and in the vicinity of the transmitter. The reception statistics show that the amplified and triggered signals were received in the conjugate region for ~25% of transmission cases for the 80 kW transmitter and over 50% for the 150 kW transmitter (Carpenter and Miller,



1976, 1983; Gibby, 2008; Li et al., 2015b). The observation occurrence was also optimized by following a procedure in which the transmission frequency and format was dynamically set and changed in response to observations of natural VLF emissions or excited echoes of transmitted signals (Gibby, 2008). In particular, transmissions within a few hundred Hz of a natural hiss band were observed to be favorable for triggering a magnetospheric response. This underlines both signal amplitude and specific frequency as important parameters in active magnetospheric whistler mode probing. Sometimes the band of frequencies over which growth occurs may be only a few hundred Hz wide (Helliwell, 1988). Signals from Siple Station were also observed on numerous spacecraft (Inan et al., 1977; Bell et al., 1981; Rastani et al., 1985). The signal amplitudes observed on the spacecraft varied from 0.01 to 0.5 pT with the strongest signals received outside the plasmopause and/or after crossing the magnetic equatorial plane. Amplitudes prior to crossing the magnetic equator were lower and in the range of 0.01 to 0.05 pT (Sonwalkar et al., 1984; Rastani et al., 1985; Sonwalkar and Inan, 1986).

The unique richness of the observations from Siple Station, which we describe in more detail in the subsequent section, motivated a wide range of theoretical studies of non-linear whistler mode wave particle interactions and the triggering of new emissions in the 1970s and 1980s (Sudan and Ott, 1971; Karpman et al., 1974, 1975; Nunn, 1974; Roux and Pellat, 1978; Vomvoridis and Denavit, 1979; Matsumoto et al., 1980). Other active experiments during this time yielded less data but confirmed the resulting effects of excited wave-particle interactions. Injection with temporary balloon transmitters (Dowden et al., 1978) or observations of pulsed VLF transmissions for maritime navigation (Tanaka et al., 1987) were also pursued. However, the former were limited by their temporary nature and the latter did not have favorable frequency and location to regularly excite the richer non-linear behavior.

Funding for Siple Station station ended in 1989 at which time the station was abandoned. A concise history of Siple Station operation can be found in Chapter 2 of the thesis by Gibby (2008); a more detailed history of experiments and operations during this period has been provided by Carpenter (2016). Data from the

Siple Station experiment originally recorded on magnetic tape has been digitized and is the subject of continued investigations (Li et al., 2014, 2015a,b; Costabile et al., 2017)

1.2. High-Frequency Active Auroral Research Program (HAARP)

The construction of the High Frequency Active Auroral Research Program (HAARP) ionospheric facility in Gakona, Alaska (62.4° N, 145.2° W) in the 1990s opened new opportunities for dedicated transmissions for magnetospheric wave injection. The main instrument of the facility is the ionospheric heater, which, upon its final completion, could radiate 3.6 MW in a wide band from 2.75 to 9 MHz, making it both the most powerful and versatile HF heater in the world. Unlike Siple Station which radiated ELF/VLF frequencies directly from a conventional antenna, the HAARP heater had the ability to generate ELF/VLF by modulating overhead natural ionospheric currents. The concept of using an ionospheric heating facility to generate ELF/VLF waves by modulating the ionospheric electrojet had been illustrated at the Tromsø facility in Norway during the 1980s (Stubbe et al., 1982) and also earlier in the Soviet Union (Getmantsev et al., 1974). However, it was initially not clear whether such a technique would be effective at HAARP since the latitude was lower than the Tromsø facility and the auroral electrojet was therefore expected to be less prominent. To the surprise of some, the first experiments of modulating the electrojet over HAARP were a huge success (Milikh et al., 1999). Even with the initial version of the heater with only 960 kW of power, ELF/VLF signals were clearly observed at a receiving station 36 km away and the facility proved effective in probing the magnetosphere (Inan et al., 2004).

The location of the HAARP ionospheric heating facility was determined by the availability of an existing military site that was originally intended to be an over-the horizon radar installation. One of the consequences of this location was that the magnetic conjugate point of the facility was in the southern Pacific Ocean about 1,000 km from the coast of New Zealand and 500 km from the nearest land of Campbell Island. An ambitious engineering effort was made to deploy autonomous receivers on buoy platforms (Cole et al., 2005) that would

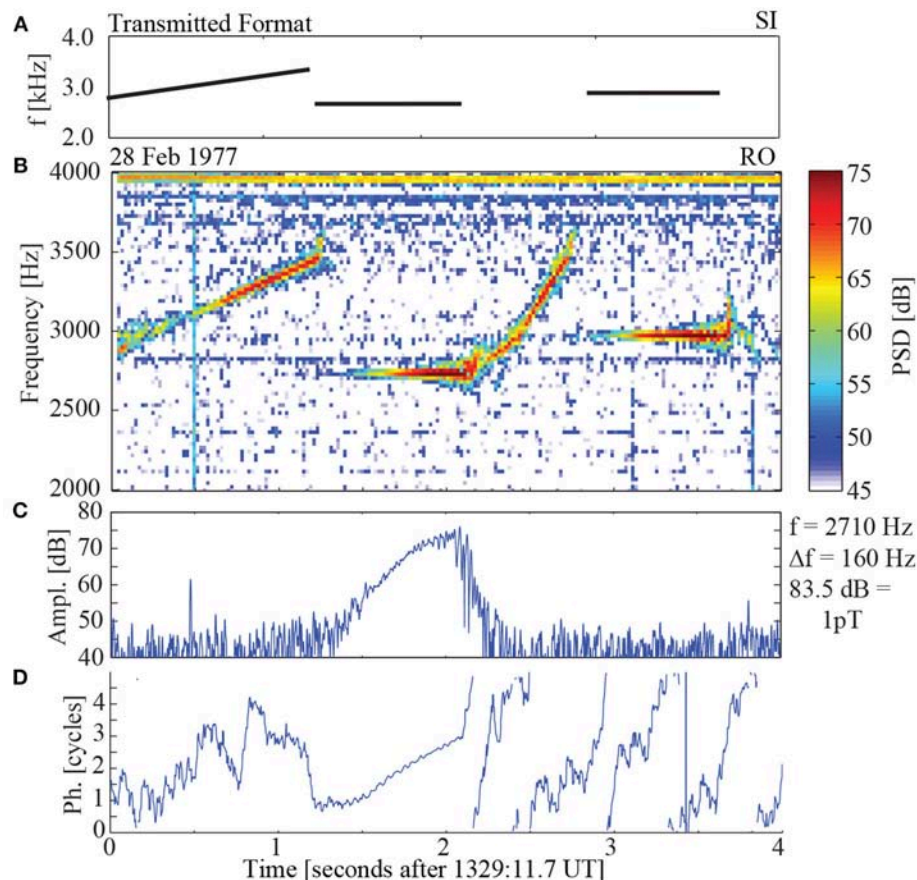


FIGURE 2 | Growth and triggering characteristics of the instability from the Siple Station experiment seen in the transmission of 2.71 kHz pulse **(A)** The time frequency format transmitted from Siple Station. **(B)** The time-frequency spectrogram of the VLF data received at the conjugate point. **(C)** A narrowband amplitude plot, centered on 2.71 kHz, with a 160 Hz bandwidth. **(D)** The phase of the narrowband data in **(C)**. Adapted from Figure 2.3 of Gibby (2008).

transmit recorded data via Iridium satellite modem. Shipborne VLF receivers were also used to make conjugate observations (Golkowski et al., 2008; Carpenter, 2016, section 5.3). The two autonomous buoy receivers deployed did not operate as long as had been initially planned but both yielded observations of one hop and higher order echoes. Likewise, almost every ship borne observation during a HAARP campaign also yielded evidence of direct whistler mode triggering by the HAARP facility. When receivers were not available in the conjugate point, a network of receivers near the HAARP facility were used to observe two hop echoes (Golkowski, 2009).

Wave injection experiments at HAARP leveraged the experience gathered during Siple Station operations. Campaigns were typically run for 1–2 weeks with ~8 h of transmissions a day. The years 2007–2008 saw a large number of campaigns dedicated to wave injection studies. The magnetospheric response to the transmissions was monitored with local receivers, which would create spectrograms in near-real time and post to a website for viewing. Changes in transmission format could be made within a minute or two by communication with the facility operator. As with Siple Station, changing the transmission frequency

and the frequency-time format would often have a significant effect on the presence and strength of magnetospheric echoes observed. HAARP ELF/VLF signals were regularly observed on the DEMETER spacecraft at 700 km altitude (Platino et al., 2006; Piddychiy et al., 2008) and also by the CLUSTER spacecraft (Platino et al., 2004). HAARP induced one-hop echoes were observed on DEMETER in the conjugate point (Golkowski et al., 2011). Additional relevant reports on HAARP wave injection include work by Golkowski et al. (2009) and Streltsov et al. (2010). A broader review of research efforts at HAARP additionally encapsulating HF wave interactions in the ionosphere has recently been compiled by Streltsov et al. (2018).

A key difference between the HAARP ELF/VLF transmissions and the Siple Station transmitter is total radiated power. As mentioned above, the Siple Station transmitter would radiate on the order of 1 kW or more of ELF/VLF power. The ELF/VLF generation capability of HAARP is variable as it depends on the overhead electrojet current intensity and the lower ionosphere profile in a complicated way (Jin et al., 2011). The radiated ELF/VLF power is also harder to quantify and estimates using both ground and space observations range from less than

1 W to a maximum of 200 W on rare select days of optimal conditions (Platino et al., 2006; Moore et al., 2007; Cohen et al., 2011; Cohen and Golkowski, 2013). Despite its lower power, the HAARP facility had the advantage of greater transmission bandwidth over spans of up to 10 kHz and ability to synthesize complex transmissions.

2. FEATURES OF OBSERVATIONS

The most characteristic repeatable feature of the observations and emblematic of the non-linear nature of the phenomena is the temporal growth in amplitude of a signal observed at a stationary receiver in the conjugate point that results from the transmission of a constant amplitude pulse. Examples of this canonical behavior from both the Siple and HAARP experiments are shown in **Figures 2, 3** respectively. The growth phase typically lasts on the order of less than a second and subsequently the amplitude saturates. During the exponential growth phase the observed frequency remains within ~ 10 – 15 Hz of the transmitted frequency, but phase is also exponentially advancing. At saturation a free running emission commences. This free running emission typically increases in frequency and is called a riser. Frequency fallers that decrease in frequency and “hook” emissions that reverse in frequency change are also observed as shown in **Figure 4**. A remarkable feature of the free running emission is that even though it swings through a wide range of frequencies of several kHz, its instantaneous bandwidth is typically restricted to less than 10 – 15 Hz. The free running emissions triggered by injection experiments are identical to features of chorus waves observed on spacecraft and on the ground. A recent analysis of two hop echoes from the HAARP experiment concurrent with chorus risers has been presented by Hosseini et al. (2017) and illustrates how the frequency sweep rate of both types of emissions shows similar evolution.

Whether or not the triggered free running emission is a riser, faller, or hook is variable. Helliwell and Katsufurakis (1974) show clear change of fallers to risers when transmitted pulse duration is changed. Fallers are generated by short pulses up to 250 ms in duration and risers are generated by longer pulses 300–400 ms long. A more comprehensive statistical study of Siple Station observations by Li et al. (2015b) confirms that shorter pulses are more likely to trigger fallers while longer pulses trigger risers. The simplest theory of the free running emission frequency change is that it is created by counterstreaming energetic electrons that are initially in forced resonance with the wave and then exit the wave field and revert back to adiabatic motion. If those electrons retain an element of phase coherence after they exit the wave field they will radiate either a falling or rising emission depending on if the reversion to adiabatic motion takes place before or after the equator (toward lower or higher gyrofrequency). This model was put forth by Roux and Pellat (1978) and is enticing in its simple elegance. However, it is noted that electrons no longer under the influence of the wave will quickly mix in gyrophase and not radiate coherently, so the distance over which this mechanism radiates would have to be small. Other more complicated theories of risers vs. fallers have these emissions

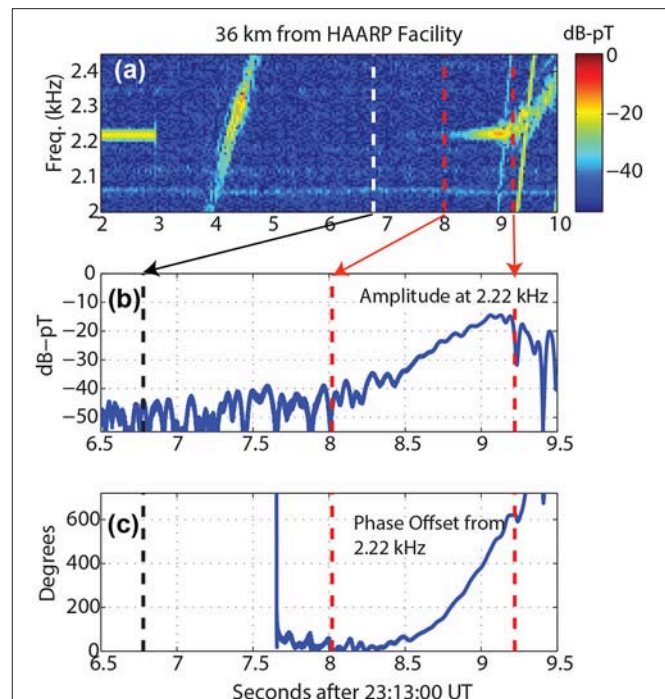
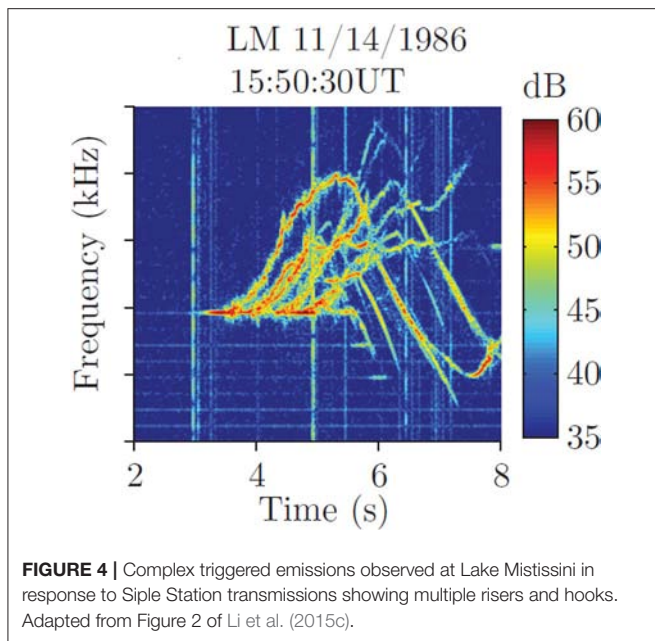


FIGURE 3 | Received two hop echo with triggered emissions due to a monochromatic transmitted pulse from HAARP at 2.22 kHz. **(a)** Spectrogram of transmitted and received signal, **(b)** temporal increase in signal amplitude at the transmitted frequency, **(c)** Exponential phase advance that accompanies the temporal amplitude increase. Features of amplitude and phase dynamics similar to those shown in **Figure 2** for a one hop echo from the Siple Station experiment. The emission at 4 s in the record is from an earlier transmitted frequency-time ramp. Adapted from Figure 3 of Golkowski et al. (2010).

being radiated by particles remaining in forced resonance with the wave but on different sides of the equator (Nunn and Omura, 2012). In either case, the magnitude and position of wave amplitude spatial gradients along the field aligned propagation path is seen as a key parameter.

2.1. Threshold for Non-linear Growth and Triggering

There is a threshold for excitation of the non-linear growth but is relatively low, on the order of a 1 W of ELF/VLF radiated power as evidenced by power stepping studies at Siple Station (Helliwell et al., 1980) and the fact that the HAARP facility was able to excite the phenomena at all given the power levels described above. In the Siple experiment, observations showing only linear growth of transmitted signals without the non-linear features were obtained (Paschal and Helliwell, 1984). For wave growth in the linear regime the echo of a transmitted single frequency constant amplitude pulse observed in the conjugate region does not show temporal amplitude change since each part of the pulse is amplified the same amount. Linear growth rates can be calculated directly from the anisotropy (see section 3.2) and flux of the energetic electron distribution (Kennel and Petschek, 1966). In the literature on Siple Station observations, linear growth is often described as “spatial” growth. Golkowski



et al. (2010) show that the exponential growth duration and also the final saturation amplitude are surprisingly similar even if the transmitted input amplitude is decreased by 13 dB. What is different in those cases is that for the weaker input the non-linear temporal growth phase occurs later and the free running emission triggered is a relatively steeper frequency riser and not a hook or faller as it is for the higher amplitude input (Gołkowski et al., 2010). The maximum saturation amplitude achieved is therefore a function of magnetospheric plasma and not the input signal amplitude. The time delay of the temporal growth and saturation can be understood as the lower input amplitude requiring more time to grow in the linear regime before the non-linear growth threshold is breached. This also means that the spatial gradients of wave amplitude along the interaction region would be in different places along the field line and thus the counterstreaming electron exiting the interaction region would exhibit different rates of gyrofrequency change when they revert back to adiabatic motion. This latter behavior can explain the riser vs. faller difference in features. In general, the non-linear temporal growth rates observed are in the range of 3 to 270 dB/s with a median growth rate of 68 dB/s (Li et al., 2015c). The observed peak echo amplitudes observed on the ground for the HAARP experiment were in the range of 0.01–1 pT. The simultaneous occurrence of both linear and non-linear wave growth makes it challenging to estimate the wave amplitude in the magnetospheric interaction region at a specific point in time and space (Gołkowski et al., 2008).

2.2. Suppression, Sidebands, and Entrainment

Other important features of the observations include suppression of growth by signals adjacent in frequency, the generation of sidebands, and entrainment. Two waves with a frequency spacing of ~ 5 Hz or less behave as a single wave and waves with

frequency spacing greater than ~ 120 Hz generate independent magnetospheric responses. Between these values, the response is suppressed relative to the independent response, with minimum response at a frequency spacing of ~ 20 Hz. The suppression occurs almost instantaneously (in less than < 10 ms) and is up to 15 dB. This suppression has been explained as stemming from the disruption of the coherent nature of a single frequency signal. Experiments testing the limits of the coherence bandwidth for triggering were performed both at Siple Station and HAARP. At Siple Station hiss like signals of band limited noise were created by modulating the frequency of the carrier. It was found that rising emissions were triggered for bandwidths less than 60 Hz but not for bandwidth at 100 Hz or greater (Helliwell et al., 1986). At HAARP, synthetic band limited Gaussian noise of instantaneous bandwidth of 10, 30, and 100 Hz was modulated onto a ELF/VLF carrier frequency. When the instantaneous bandwidth was 30 Hz or below, magnetospheric amplification and triggering was observed, when it was 100 Hz no amplification was observed (Gołkowski et al., 2011). These results suggest that hiss emissions can trigger or evolve to discrete chorus like emissions but only if a minimum level of coherence or maximum bandwidth is achieved. In this context, observations made by Hosseini et al. (2017) show a band of hiss narrowing in bandwidth before the hiss emissions evolve to chorus emissions.

Entrainment is a multi-frequency interaction in which an injected signal captures a free running emission and controls its frequency (Helliwell and Katsufakis, 1974; Gibby, 2008; Gołkowski et al., 2008). An example of free running emissions being successively entrained by a series of transmitted pulses decreasing and then increasing in frequency is shown in Figure 5. This phenomenon shows how the hot plasma distribution that is radiating the free running emission can be directly modified in a very deterministic manner. In a broader context, modification of the frequency content of chorus waves may be possible if controlled wave power can be injected at the appropriate place and time.

Sidebands occur when one or more quasi-constant frequency components appear within less than 100 Hz of a monochromatic input wave. The term sidebands originates from the overall similarity to modulated radio communications. Sidebands appear rarely and when the observed carrier wave is strong, but there is no simple relationship between carrier amplitude and sideband amplitude. Sideband amplitude may be symmetrical or asymmetrical about the carrier, and in the asymmetrical case it is usually the upper sideband that is stronger. Sideband amplitude is usually 10 dB or more below the carrier amplitude, but sometimes it can exceed the carrier amplitude (Park, 1981). Costabile et al. (2017) performed relative phase analysis of sidebands from the Siple experiment and discuss theories of sideband generation.

2.3. Effect of Geomagnetic Conditions and Transmitted Frequency

Both the Siple Station and HAARP experiments found that observations of magnetospheric echoes were most likely after 2–3 days of quieting geomagnetic conditions following a

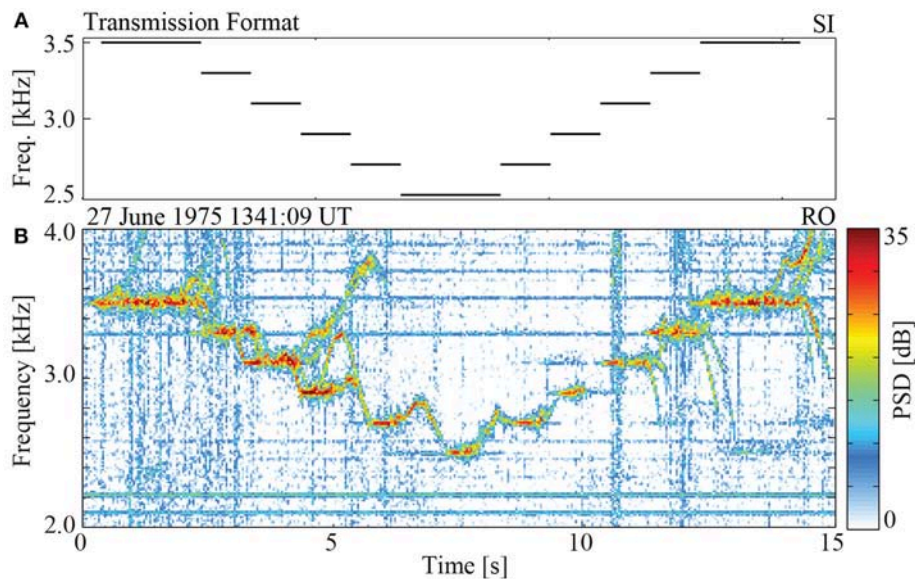


FIGURE 5 | A case of entrainment from the Siple Station experiment received at Roberval on 27 June 1975 at 13:41:09 UT, in which a “stair-step” transmission from Siple starting at 3.5 kHz triggered a free-running plasma emission whose time-frequency characteristics are controlled by the subsequent constant frequency transmissions. **(A)** Time-frequency format transmitted by Siple. **(B)** Received signal at Roberval (magnetic conjugate point), showing characteristics of entrainment. Adapted from Figure 2.7 of Gibby (2008).

magnetospheric disturbance (Carpenter and Bao, 1983; Helliwell, 1988; Gołkowski et al., 2011; Li et al., 2015c). Although highly disturbed conditions and the associated free energy of high radiation belt fluxes can seem favorable for triggering of non-linear phenomena, the requirement for a stable ducted path for propagation of a ground injected signal to the equatorial interaction region appears to be a dominant factor. The latter condition is known to be associated with quieting conditions. For the Siple Station experiments the vast majority of signal receptions occurred when the transmitted frequency was between 0.2 and 0.5 of the equatorial gyrofrequency (~ 6 kHz for $L = 4.2$). For the HAARP experiment, the observed echoes were always below half the equatorial gyrofrequency (~ 3.8 kHz for $L = 4.9$). A common explanation put forth for very few observations above the equatorial half gyrofrequency has been that guiding of waves in density enhancement ducts is limited by this upper cutoff (Smith, 1960) (in depletion ducts, propagation above the half the gyrofrequency is possible.) However, recent work suggests that for typical pitch angle anisotropy values found in the magnetosphere, linear growth also exhibits an upper frequency cutoff. For anisotropy values of 1, the cutoff is exactly at half the gyrofrequency (Hosseini et al., 2019). Additional theories posit that the observed half-gyrofrequency cutoff may be due to non-linear damping of longitudinal components due to quasi-parallel propagation (Omura et al., 2009; Yagitani et al., 2014).

2.4. Terminology

The richness of the observations has caused a number of terms to be used to describe the phenomena and this can lead to confusion. In the literature describing Siple Station results the term “coherent wave instability”(CWI) is commonly used to

emphasize that the injected wave needs to be phase coherent over a minimum temporal duration for the non-linear interaction to take place. This minimum duration requirement was also seen in the “dot-dash” anomaly of the early observations. The prominent stages of the single frequency excitation, namely temporal growth near the transmitted frequency followed by saturation, and a “free” running riser or faller are also often parsed with the terms “echo,” “triggering wave,” or “embryo emission” for the initial part (Dowden et al., 1978) and “triggered emission” for the latter free running component. On the other hand, theoretical publications by authors removed from the active experiments tend to use the term “VLF triggered emissions” to describe all the phenomena. Due to the common physics between wave injection results and natural chorus waves, a broader term of magnetospheric non-linear cyclotron growth is perhaps the most appropriate.

3. THEORY OF CYCLOTRON RESONANCE AND WAVE AMPLIFICATION

The fundamental physical environment of cyclotron wave-particle interactions in the magnetosphere is reasonably well-understood (Gendrin, 1975; Omura et al., 1991; Thorne, 2010). Specifically, in the region of the plasmasphere $2 < L < 6$, the Earth’s magnetic field retains an approximately dipole shape. Additionally, a population of low energy $1 \text{ eV} < E < 10 \text{ eV}$ but relatively dense $10 \text{ cm}^{-3} < N_c < 5,000 \text{ cm}^{-3}$ electrons permeate the background which results in a magnetized plasma environment. The plasmasphere thus supports the propagation of several plasma wave modes of which, as discussed previously,

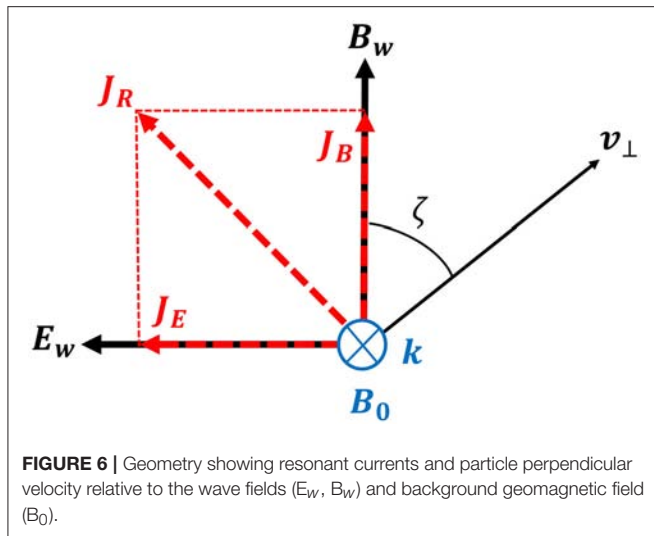


FIGURE 6 | Geometry showing resonant currents and particle perpendicular velocity relative to the wave fields ($\mathbf{E}_w$, $\mathbf{B}_w$) and background geomagnetic field ($\mathbf{B}_0$).

the whistler mode is of particular importance. Superimposed on the cold particles are the outer Van Allen radiation belts which consist of high energy electrons $1 \text{ keV} < E < 100 \text{ keV}$ that are trapped in a magnetic mirror configuration by the geomagnetic field (Walt, 2005). Since radiation belt electrons are forced into helical orbits by the background field, the particles can resonate with the circularly polarized whistler mode waves that are also propagating along the field line. Electrons that are counter streaming to the wave's propagation direction can undergo Doppler shifted cyclotron resonance, or gyro-resonance. That is, electrons that travel at the appropriate velocity will experience an approximately static wave electric field and significant energy exchange Inan (1977). This is referred to as the resonance velocity, v_r and is given by

$$v_r = \frac{\omega - \frac{\omega_c}{\gamma}}{k} \quad (1)$$

where the quantities ω and k correspond to the wave frequency and wavenumber respectively and are related by the whistler mode dispersion relation. The quantity γ is relativistic Lorentz factor and is important for ultra-relativistic particles (Omura et al., 2007). An important assumption in (1) is that the waves are assumed to propagate parallel to the magnetic field lines and all other wave modes are ignored. This is a reasonable assumption when assuming ducted propagation although some work may suggest the importance of additional wave modes as well (Bell and Ngo, 1990; Zhang et al., 1993). Recent work has shown that the results of parallel propagation are still applicable for small oblique angles of propagation (Nunn and Omura, 2015). Note, as per the convention of Omura et al. (2008), the waves propagate in the z -direction and the resonance velocity is thus negative for counter streaming electrons.

The mathematical basis of modeling wave-particle interactions is via the Vlasov-Maxwell system of equations. Specifically, the Vlasov equation (2) describes the evolution of the electron phase space density $f(\mathbf{r}, \mathbf{v})$ in a collision-free plasma.

$$\frac{\partial f}{\partial t} + \mathbf{v} \cdot \frac{\partial f}{\partial \mathbf{r}} - \frac{q}{m} (\mathbf{E}_w + \mathbf{v} \times \mathbf{B}) \cdot \frac{\partial f}{\partial \mathbf{v}} = 0 \quad (2)$$

Here, the quantities $\mathbf{r}$ and $\mathbf{v}$ correspond to the position and velocity coordinates of phase-space. E_w corresponds to be the wave electric field while $\mathbf{B}$ is the total magnetic field. The total magnetic field can be decomposed into $\mathbf{B} = \mathbf{B}_w + \mathbf{B}_0$ where $\mathbf{B}_w$ is the wave magnetic field while $\mathbf{B}_0$ represents the background geomagnetic field.

Maxwell's equations govern the evolution of the wave electric and magnetic fields and are given by (3-4),

$$\nabla \times \mathbf{E}_w = -\frac{\partial \mathbf{B}_w}{\partial t} \quad (3)$$

$$\nabla \times \mathbf{B}_w = \mu_0 (\mathbf{J}_h + \mathbf{J}_c) + \frac{1}{c^2} \frac{\partial \mathbf{E}_w}{\partial t} \quad (4)$$

The quantities $\mathbf{J}_h$ and $\mathbf{J}_c$ represents the currents due to the hot and cold plasma respectively.

Although the general theory of wave particle interactions is rather complex, analytical expressions can be derived under certain simplifying assumptions. Specifically, wave growth is typically separated into two regimes, (i) linear growth driven by temperature anisotropy and (ii) non-linear growth driven by phase-trapping of resonant particles. The two regimes are both important components of the whistler mode instability and are discussed in more detail in the following subsections.

3.1. Narrowband Field Equations and Geometry

When modeling the evolution of the electric and magnetic fields in a magneto-plasma, the wave equations can be simplified under the assumption of a narrowband modulating wavepacket. This is a reasonable assumption given the coherence of the signals observed in the data. Specifically, the expression for a circularly polarized whistler wave magnetic field propagating in the $+z$ -direction is given by

$$\mathbf{B}_w = \Re \left[(\hat{\mathbf{x}} - j\hat{\mathbf{y}}) B_w e^{j(\phi_w + \omega t - \int k dz)} \right] \quad (5)$$

where $j = \sqrt{-1}$. The term $\omega t - \int k dz$ in the argument of the exponent corresponds to the phase variation of a monochromatic plane wave and can be thought of as a feature of the injected carrier wave. The quantity $B_w e^{j\phi_w}$ corresponds to the complex wavepacket that modulates the carrier whistler wave. Under the slowly-varying or narrowband assumption (Nunn, 1974) the evolution equations for the amplitude and phase of the modulating wavepacket is given by (6)-(7),

$$\left(\frac{\partial}{\partial t} + v_g \frac{\partial}{\partial z} \right) B_w = -\frac{\mu_0 v_g}{2} J_E \quad (6)$$

$$\left(\frac{\partial}{\partial t} + v_g \frac{\partial}{\partial z} \right) \phi_w = -\frac{\mu_0 v_g}{2} \frac{J_B}{B_w} \quad (7)$$

These narrowband wave equations describe the evolution of a wavepacket that is propagating at the group velocity of the whistler wave. Specifically, (6) shows that the wave amplitude

is driven by the component of the resonant current that is parallel to the wave electric field J_E . When J_E is negative the wave will experience growth, otherwise the wave will be damped for positive J_E . On the other hand, the wave phase (and hence frequency change) is driven by the component of the current that is anti-parallel to the wave magnetic field J_B . The geometry that shows the wave fields, the resonant currents, and the particle velocity variables are delineated in **Figure 6**. The quantity $\mathbf{J}_R$ represents the resonant current due to the hot plasma. (J_E and J_B are orthogonal components of $\mathbf{J}_R$) The variables $v_\perp$ and ζ correspond to each particle's velocity perpendicular to $\mathbf{B}_0$ and the gyrophase angle respectively.

Equations (6–7) have been derived independently by several authors (Karpman et al., 1974; Nunn, 1974; Rathmann et al., 1978; Omura and Matsumoto, 1982; Trakhtengerts, 1995) and are believed to adequately describe the wave fields for the whistler mode instability.

3.2. Linear Theory

Under the assumption of small amplitude waves and small perturbations to the initial particle distribution, the Vlasov-Maxwell system can be linearized and a closed form expression can be derived accordingly (Kennel and Petschek, 1966). Following the method of Golkowski and Gibby (2017), the resonant currents under the linearized assumption are given by (8, 9),

$$J_E = -\frac{2}{\mu_0 v_g} \gamma_L B_w \quad (8)$$

$$J_B = 0 \quad (9)$$

where γ_L is the well-known linear growth rate (Kennel and Petschek, 1966) and is given by the expression (10),

$$\gamma_L = \pi \frac{\omega_c}{N_c} \left(1 - \frac{\omega}{\omega_c}\right)^2 |v_r| \left[A - \frac{\omega}{\omega_c - \omega}\right] \eta \quad (10)$$

N_c represents the cold plasma density, while η and A correspond to the particle resonant flux and anisotropy respectively. As can be seen in (10), the sign of the linear growth rate is dictated by the value A . In the case of a Bi-Maxwellian velocity distribution, the expression for the anisotropy simplifies to

$$A = \frac{T_\perp}{T_\parallel} - 1. \quad (11)$$

The quantities $T_\perp$ and $T_\parallel$ correspond to the electron temperatures in the directions that are perpendicular and parallel to geomagnetic field respectively. Thus, if the electron distribution has sufficient temperature anisotropy ($A > \frac{\omega}{\omega - \omega_c}$) the radiation belt velocity distribution is unstable and whistler waves can be amplified. An interesting consequence of (8, 9) is that only J_E is non-zero while J_B is identically zero under the linearized model. Thus, linear theory predicts amplification and no frequency change of the injected whistler wave. However, observations as discussed in section 1 show frequency changes

as defining features of the non-linear instability. As such, linear theory only describes the initial process of wave amplification and cannot be used to model the instability in its entirety.

3.3. Non-linear Theory

Once the magnetic field of the wave becomes sufficiently large, linear theory does not adequately model the whistler mode instability. Accurately understanding the dynamics of resonant particles is required to correctly describe the non-linear aspect of the problem.

The dynamics of an energetic electron in a monochromatic whistler mode wave field is in general governed by the Lorentz force. Although several authors have analyzed the equations of motion with different approaches, this section will review the simplified equations that are generally accepted to be the most relevant. The equations of motion can be simplified by neglecting the transverse spatial motion of electrons and by only considering spatial variation along the field line coordinate, z . Additionally, by using a cylindrical coordinate system in velocity and only considering the dynamics of near-resonant particles (Omura et al., 1991), the equations of motion can be written as (12, 13),

$$\frac{d\zeta}{dt} = \theta \quad (12)$$

$$\frac{d\theta}{dt} = \omega_{tr}^2 (\sin \zeta + S) \quad (13)$$

Here, the variable $\theta = k(v_\parallel - v_r)$ represents a normalized change of the electron's parallel velocity from resonance. The quantity $\omega_{tr} = \sqrt{\frac{qkv_\perp B_w}{m}}$ is known as the trapping frequency. The quantity S is called the collective inhomogeneity factor or the “S-parameter” and is given by

$$S = -\frac{1}{\omega_{tr}^2} \left[\left(\frac{kv_\perp^2}{2\omega_c} + \frac{3}{2} |v_r| \right) \frac{\partial \omega_c}{\partial z} + \frac{2\omega + \omega_c}{\omega} \frac{d\omega}{dt} \right]. \quad (14)$$

The S-parameter quantifies the effect of background inhomogeneity as well as the frequency sweep rate as observed by the particle ($\frac{d\omega}{dt}$). It is worth noting that several authors have derived (12), (13), and (14) with different notation over the past several decades (Dysthe, 1971; Nunn, 1974; Matsumoto and Omura, 1981; Trakhtengerts and Rycroft, 2008). Differentiating (12) with respect to time and plugging into (13), results in a non-linear ordinary differential equation given by

$$\frac{d^2 \zeta}{dt^2} = \omega_{tr}^2 (\sin \zeta + S). \quad (15)$$

Equation (15) represents a forced pendulum equation where the forcing term is proportional to S . For $S = 0$, (15) is identical to the conventional pendulum equation and the particle will oscillate around $\zeta = \pi$ at the trapping frequency ω_{tr} in a manner similar to which a pendulum oscillates in a constant gravitational field. For values in the range $-1 < S < 1$, the central phase angle around which the particle oscillates is moved to $\zeta_0 = -\arcsin(S)$. For $|S| > 1$ particles are not trapped and do not remain in

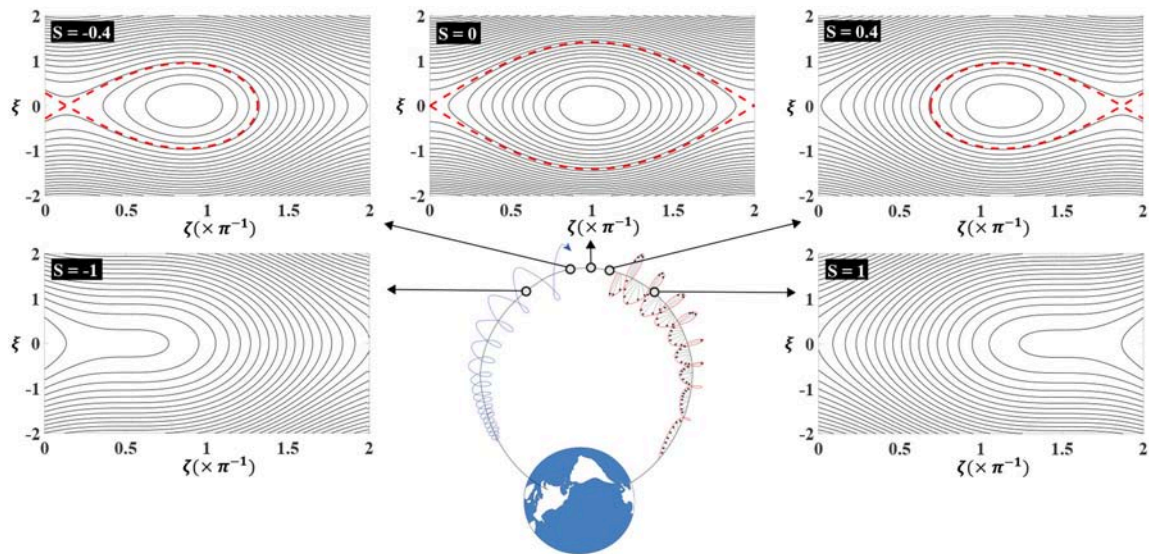


FIGURE 7 | Local phase-space trajectories at different locations along the field line for a monochromatic and constant amplitude signal. The separatrix is shown by the red contours. A zero value on the vertical axis represents the resonance velocity given in Equation (1). The horizontal axis is gyrophase as defined in Figure 6.

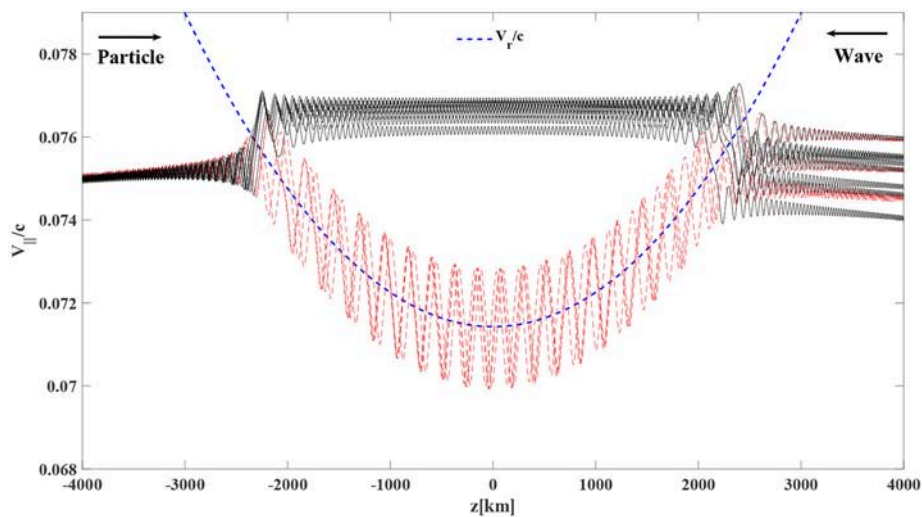


FIGURE 8 | Trajectories of trapped (red) and untrapped (black) electrons with uniformly distributed initial gyro-phase. The wave is assumed to be constant amplitude and monochromatic.

resonance with the wave. For a dipole geomagnetic field that is typically accurate in the plasmasphere and $\frac{d\omega}{dt} = 0$, $S = 0$ only at the magnetic equator. For a distorted geomagnetic field geometry so called “minimum B” pockets can occur off the equator and also enhance particle phase trapping even for low amplitude waves. In this context it is worth noting that chorus waves are observed to be primarily generated at the equator (Santolik and Gurnett, 2003) or in such minimum B pockets (Tsurutani and Smith, 1977).

Of particular importance is the formation of a wave-induced trap in phase-space (Omura et al., 2008). Figure 8 shows twelve test particle trajectories with trapped resonant particles (red)

and untrapped resonant particles (black) for an assumed dipole geomagnetic field. All particles start with the same value of $v_{\parallel}$ and $v_{\perp}$ as well as the same initial position. The particles are uniformly distributed in gyrophase, and as shown in Figure 8, the untrapped particles are deflected when they come into resonance with the wave. On the other hand, the trapped particles are forced to stay in resonance with the wave over thousands of kilometers after which they are released from the trap. Whether or not a specific particle is trapped depends on the initial gyrophase angle when the particle goes into resonance with the wave. Phase-trapping of particles is believed to be a vital component of the non-linear instability (Dysthe, 1971; Matsumoto et al., 1980;

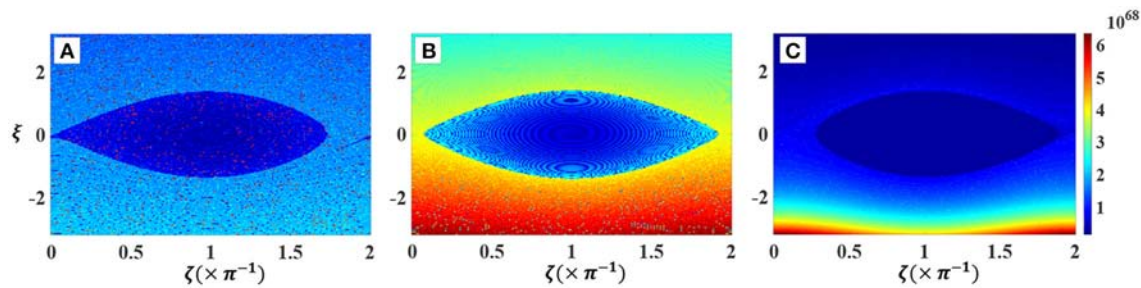


FIGURE 9 | Formation of phase-space electron hole from test particle trajectories. (A) upstream location, (B) equator, and (C) downstream location relative to the wave propagation direction.

Harid, 2015). Specifically, trapped particles deviate significantly from their adiabatic trajectories which in turn appreciably alter the distribution function.

The structure of the phase-space trap depends on location along the field line and **Figure 7** shows the shape of the phase-space trap at several positions (for a monochromatic whistler mode wave). The variable ζ on the vertical axis is defined by $\frac{\theta}{\sqrt{2\omega}}$ and is essentially a normalized deviation of $v_{\parallel}$ from resonance. The trapped trajectories correspond to closed curves in phase-space while the untrapped particles follow open curves. The trapped and untrapped electron populations are separated by a boundary known as a separatrix and are shown by the red contours. For a monochromatic signal, the phase-space trap can only exist in a narrow range around the magnetic equator after which the trap disintegrates and the mirror force dominates over the non-linear effects of the wave.

Since trapped particles that are downstream of the wave are forced to remain in resonance for a long period of time, by virtue of Liouville's theorem the trapped particles drag the downstream value of the distribution function to locations that are upstream (Here upstream and downstream are defined relative to the wave propagation direction). Within a few trapping periods, the density inside the trap will be approximately constant (i.e. phase-mixed) while the region outside the trap will be close to the unperturbed velocity distribution. As shown in **Figure 8**, electrons that are trapped downstream will start at high value of $v_{\parallel}$ and follow the resonance velocity curve to a lower values of $v_{\parallel}$ at the equator. Since the initial velocity distribution typically has a lower value at higher particle velocities, the density inside the trap at the equator will be much lower than the surrounding regions of phase-space. This results in what is known as an "electron hole" in phase-space.

By running test particle trajectories backwards in time and employing Liouville's theorem, the distribution function can be reconstructed in high resolution (Nunn, 2012; Harid et al., 2014a). **Figure 9** clearly shows the electron hole in for three different locations along the field line (upstream, equator, and downstream of the wave) in the presence of a monochromatic and constant amplitude signal. As shown, the electron hole is well-defined and has an approximately constant density inside the phase-space trapping region. It is worth noting that for higher pitch angles or short pulses the opposite can occur and

an "electron hill" can be formed as well (Hikishima and Omura, 2012; Nunn and Omura, 2012).

The formation of an approximately constant density across the phase-space trap allows for semi-analytical calculation of the resonant currents (Omura et al., 2008, 2009; Summers et al., 2012). Omura et al. (2009) and Cully et al. (2011) used such expressions along with further assumptions to estimate frequency sweep rates of chorus emissions. Costabile et al. (2017) used such expressions to investigate sideband formation. Although these simplifications have been validated against simulations (Kato and Omura, 2016) and satellite observations (Cully et al., 2011) of chorus emissions, it does not entirely describe the complex dynamics of triggered VLF emissions. This is partly because triggered emissions are induced by a coherent seed wave while chorus waves are generated from amplification of background noise that is maximized at the equator (Foust, 2012). The formation of triggered waves, however, may occur at a point along the field line that is offset from the equator resulting in falling tones, hook like emissions, and other complex frequency-time relations (Smith and Nunn, 1998). Thus, the theoretical framework requires some extension to handle the rich variation that is observed in data (Helliwell, 1965; Li et al., 2015b).

Several theoretical features of the whistler instability, particularly the basis of amplification, has been well studied over the past several decades. However, many important features are yet to be properly understood from a theoretical point of view. These include complex variability of rising, falling, and hook-like emissions. The interaction between multiple waves that are closely spaced in frequency (and thus have overlapping traps), can lead to what is known as the coherence bandwidth effect and has not been considered in rigorous detail. Additionally features such as the entraining of one signal onto another has yet to be accurately described from fundamental physics. Effects of additional plasma modes or three dimensional aspects of the real physical scenario have often been neglected and some research suggests that there may be important physical phenomena that is yet to be captured (Omura and Matsumoto, 1987; Bell and Ngo, 1990; Ke et al., 2017). It is likely that theoretical insight will be gained via numerical simulations, especially in the current era of high performance computing paradigms.

4. SIMULATION METHODS AND RESULTS

The inherent non-linearity associated with non-linear whistler phenomena is largely analytically intractable and thus numerical simulations are the primary means of approaching the problem. From a theoretical point of view, the cyclotron instability can be modeled using the Vlasov-Maxwell system of equations. Although the problem is well-defined in theory, the multi-scale aspect of the problem requires care when developing self-consistent simulations. Specifically, the electron resonant velocity in the presence of a monochromatic wave varies significantly along the field line due to the spatial variation of the geomagnetically field. On the other hand, the size of the trap in phase-space ($v_{tr} = \frac{\omega_{tr}}{k}$) is much smaller than the typical values of the resonance velocity. Additionally, adiabatic motion of particles that are outside the phase-space trap cover a very large range of velocities due to the geomagnetic mirror geometry. Thus, the simulations must resolve the trap with enough detail to discern non-linear dynamics while the range of particle velocities must be large enough to encompass all the resonance velocities and adiabatic trajectories that constantly fall into resonance. These complications make simulations difficult and simplifying assumptions are often needed for computational feasibility at the expense of ignoring certain physical effects. For this reason, several computer models have been developed over the past five decades, each of which have various strengths and weaknesses. Although several authors have considered test particle dynamics in the presence of whistler mode waves (Inan, 1977; Albert, 2002; Tao et al., 2012; Albert et al., 2013), only models that self-consistently account for wave amplification are considered in this review. Reviewing the historical development of various numerical models helps provide a good understanding of the major contributions so far as well as the unanswered questions that are relevant for future work.

4.1. Modeling History and Results

Self-consistent computer simulations of the non-linear whistler mode instability in the magnetosphere have been utilized extensively since the late 1960s. Only certain major works are described in this review to illustrate the historical progression of computational techniques. Thus, the ensuing discussion is by no means exhaustive and is meant to provide a high level view of simulation results over the past 50 years.

The earliest simulations can be traced back to Helliwell and Crystal (1973) where the authors considered radiation due to a monoenergetic stream of resonant sheets of phase-bunched electrons. Although the model predicted wave-growth, the important effect of the geomagnetic field inhomogeneity was ignored as well as the changing frequency of the wave. Additionally, the model did not consider a realistic initial electron distribution function which plays an important role in the wave-particle interactions process. Nunn (1974) developed a hybrid code where the cold particle population was modeled via a fluid equation while the resonant currents were assumed to be dominated by stably trapped particles and the wave equations were approximated under a narrowband assumption. Additionally, a phenomenological damping term was included

to account for effects such as Landau damping and leakage from a duct. The model did produce non-linear amplification and demonstrated the formation of currents due to resonant interactions, however, frequency change was not readily observed in the simulations. This was primarily due to the fact the range of particle velocities in the simulations were quite close to the local resonance velocity in the presence of a monochromatic wave-packet, thus velocities corresponding to resonance at new frequencies were inherently ignored. Denavit and Sudan (1975) utilized a full particle simulation where both the cold and hot plasma are modeled with a large number of macro-particles. Just as in Nunn (1974), the waves were treated with a narrowband assumption for computational simplicity. The authors showed that the model did produce non-linear amplification for an unstable plasma. Additionally elongation and frequency change of the wavepacket was observed due to phase correlations of detrapped resonant electrons.

The work of Denavit and Sudan (1975) was one of the first full-particle simulations of the whistler-mode instability. Similarly, Vomvoridis and Denavit (1980) applied the long-time scale (LTS) algorithm of Rathmann et al. (1978) to the wave-particle interactions problem. The model essentially tracked particle trajectories through time and the resonant currents were determined by giving the particles a finite size in a manner similar to Denavit and Sudan (1975). Unlike the Vlasov Hybrid Simulation (described below), no phase-space grid was required for the simulations. They found that the growth could be separated into a homogeneous component, inhomogeneous untrapped component, and an inhomogeneous trapped component. Although the model elucidated features of non-linear growth due to resonant wave-particle interactions, the simulated frequency change primarily showed temporal oscillations that were in part attributed to undersampling of the particle distribution. The code utilized a narrowband assumption for the wave envelope without any filtering which may have caused further difficulty in modeling free running emissions. Additionally, the code was highly susceptible to numerical noise and oscillations due to the undersampling problems associated with particle methods. Matsumoto et al. (1980) and Omura and Matsumoto (1985) also considered full particle simulations but with a homogeneous background magnetic field. The work clearly demonstrated non-linear growth, however, significant frequency changes was not observed which further highlighted the importance of a spatially varying geomagnetic field.

One of most impactful numerical models was the development VHS (Vlasov Hybrid Simulation) code by Nunn (1990). The code included the effects of trapped and untrapped resonant particles and relied on continuously tracking particle trajectories in time while interpolating to a phase-space grid with the aid of Liouville's Theorem. The resonant currents were then calculated by appropriately integrating over phase-space. The VHS code successfully reproduced several features of triggered VLF emissions including non-linear growth, rising frequency tones, falling tones, and hook-like signals. Although the code shed light on several aspects of free-running emissions, it required artificial filtering to ensure a well-defined frequency as well as the stability of the simulations. Additionally

the narrowband assumption make it unsuitable for complex multi-frequency interactions or broadband signals. Even so, the VHS code has been successfully used to reproduce several features observed in data (Nunn et al., 1997, 2009; Smith and Nunn, 1998).

Increased computational power in the early 2000s permitted the design of higher resolution codes and renewed interest in modern simulations of the whistler mode instability. Specifically, the ability to use several million to even billions of grid points has become a practical reality with parallel computing paradigms and decreased memory costs. Additionally, simplified models with fewer grid points were capable of being run on a desktop computer within a few hours. Gibby et al. (2008) used a model similar to Nunn (1990) with a narrowband hybrid approach, however, the particle trajectories were not calculated continuously but were reset at every time step. This is known as a semi-Lagrangian method and typically results in a smoother interpolation of the particle distribution (Sonnendrücker et al., 1999). The model demonstrated saturation of coherent waves as well the beginning stages of frequency change. The window of particle velocities in the simulations were a relatively narrow sliver around resonance, thus large changes in frequency was not supported by the code. Harid et al. (2014b) also used a similar approach, however, a canonical resonant coordinate transformation was used to employ a finite difference scheme in phase-space. The number of grid points was an order of magnitude higher than used by Gibby et al. (2008) which provided high resolution features in phase-space. Specifically the clear formation of a density depletion in the phase-space trap was observed during the non-linear growth phase, which confirmed the long-standing hypothesis of the electron “phase-space hole” as a dominating feature of the wave-growth process (Omura et al., 2009).

Katoh and Omura (2006) developed the first modern hybrid-particle simulation where the cold plasma was treated as a fluid and the hot plasma was modeled with a PIC approach. The simulation used approximately 67 million particles and successfully produced rising tone triggered emissions. The distinguishing feature from previous work is that the initial distribution function was a Bi-Maxwellian where all particles (resonant and non-resonant alike) were taken into account. Additionally, the model did not utilize a narrowband assumption for the waves and Maxwell's equations were solved with a fourth order finite difference time domain (FDTD) scheme. Thus, the code can in principle be used to model narrowband wave-particle interactions, multi-wave interactions, as well as broadband injected signals. The code was later successfully utilized to simulate the generation of naturally generated chorus and hiss waves (Katoh and Omura, 2008, 2016; Omura et al., 2009) which further demonstrated the robustness and utility of the hybrid-particle model.

Hikishima et al. (2009) utilized the first full particle simulation to successfully model chorus emissions. The code was then used to model triggered emissions by injecting a constant frequency signal at the equatorial region of the simulation (Hikishima et al., 2010). The model used approximately 150 million particles and successfully simulated triggered rising tone emissions along with

the corresponding amplification of the seed and triggered waves. The large number of particles were required to overcome the noisy fluctuations associated with particle codes. The results also showed evidence of the generation of rising tone emissions at the back end of the seed signal, indicating both the importance of trapping as well as detrapping of resonant particles. Hikishima and Omura (2012) used the same particle code to run a parametric study by varying the injected wave amplitude. The authors found that either extremely small ($B_w < .2 \times 10^{-3} B_0$) or extremely large amplitudes ($B_w > 4 \times 10^{-3} B_0$) did not result in rising tone emissions, where B_0 is the value of the background magnetic field strength at the equator. Additionally, in the range of amplitudes where triggered emissions were generated, a clear formation of a hole in phase-space was observed in the simulations. This code was the first to show the formation of an electron hole while employing a broadband particle simulation without simplifying assumptions.

The work of Katoh and Omura (2008) and Hikishima and Omura (2012) demonstrates the current state-of-the-art in modeling the whistler mode instability in the magnetosphere with controlled excitation. However, all the mentioned works so far have the inherent limitation of being one-dimensional in space and three dimensional in velocity. Additionally, the simulations all consider parallel propagation with only electromagnetic plasma waves so the electrostatic components have been neglected. Ke et al. (2017) was the first published work that employed a two dimensional hybrid-PIC code to simulate the generation of chorus emissions. The results showed that chorus waves are generated close to the magnetic equator and increase in wavenormal angle during propagation to higher latitudes. Although the code has not been utilized for excitation by injected waves, the physical mechanism behind triggered emissions and chorus waves are similar. Thus, the method used by Ke et al. (2017) is a powerful approach to model higher dimensional effects of controlled wave excitations.

A summary of several numerical models that have been developed for the whistler instability in the magnetosphere is shown in Table 1.

4.2. Types of Numerical Models

Given the several codes that have been successfully utilized over the past several decades, it is useful to categorize the various self-consistent models into general types. The most general simulation requires providing a numerical approximation to the Vlasov-Maxwell system of Equations (2), (3), and (4). This type of solution is typically referred to as a fully kinetic simulation. A fully kinetic simulation treats both the cold plasma and radiation belt particles via a Vlasov approach. However, since kinetic effects of cold plasma particles are effectively negligible, a common methodology is to consider a hybrid-kinetic approach. In hybrid methods, the cold electrons are treated as a fluid while the hot electron evolution is governed by the Vlasov equation. As far as solving the Vlasov equation for the hot plasma, most solvers can be lumped into two general categories, Vlasov continuous codes (VCON) or particle codes (PIC) (Filbet et al., 2001; Gutnic et al., 2004).

TABLE 1 | Key self-consistent numerical codes used to simulate nonlinear cyclotron resonance and wave growth.

References	Solver type	Code features	Key results
Numerical models			
Helliwell and Crystal, 1973	Hybrid-PIC	1R-3V Narrowband Homogeneous Resonant Electrons	Wave growth No Frequency Change
Nunn, 1974	Hybrid-VCON	1R-3V Narrowband Includes Inhomogeneity Resonant Electrons	Wave growth Small Frequency Change
Denavit and Sudan, 1975	Hybrid-PIC	1R-3V Narrowband Includes Inhomogeneity Near-resonant Electrons	Wave growth Frequency Change
Vomvouridis and Denavit, 1980	Hybrid-PIC	1R-3V Narrowband Includes Inhomogeneity Near-resonant Electrons	Wave growth Frequency change
Matsumoto and Omura, 1985	Full-PIC	1R-3V Narrowband Homogeneous/Inhomogeneous Near-resonant Electrons	Wave growth No Frequency change
Nunn, 1993	Hybrid-VCON	1R-3V Narrowband Includes Inhomogeneity Near-resonant Electrons	Wave growth Triggered Emissions
Gibby et al., 2008	Hybrid-VCON	1R-3V Narrowband Includes Inhomogeneity Near-resonant Electrons	Wave growth Short Triggered Emissions
Harid et al., 2014b	Hybrid-VCON	1R-3V Narrowband Includes Inhomogeneity Near-resonant Electrons	Wave growth Short Triggered Emissions
Kato and Omura, 2006	Hybrid-PIC	1R-3V Broadband Includes Inhomogeneity All Electrons	Wave growth Triggered Emissions
Hikishima et al., 2009	Full PIC	1R-3V Broadband Includes Inhomogeneity All Electrons	Wave growth Triggered Emissions
Ke et al., 2017	Hybrid-PIC	2R-3V Broadband Includes Inhomogeneity All Electrons	Wave growth Triggered Emissions 2D Propagation Effects

4.2.1. VCON Methods

The distinguish feature of VCON codes is that they rely on creating a grid in phase-space and determining the value of the distribution function on these grid points. The generation of a phase-space grid is often referred to as an Eulerian method. The currents are then calculated by appropriately integrating over phase-space. There are several possible techniques that utilize a phase-space grid, however, the few methods that have been applied to the whistler-mode instability will be discussed.

The method employed by Nunn (1990) can be considered a semi-Eulerian or semi-Lagrangian method. For simplicity

of presentation, we consider a two-dimensional phase-space (z, v) , however the analysis naturally translates over to higher dimensions. In this technique, the initial distribution function $f(z, v)$ is first initialized on a grid in phase-space of size $N_z \times N_v$ where each grid point has the coordinates (z_n, v_m) and volume $\Delta z \Delta v$ for $n = 1, 2, \dots, N_z$ and $m = 1, 2, \dots, N_v$. Each cell on the grid can then be thought of as a “super-particle” with density $f_{nm} = f(z_n, v_m)$ and thus charge $Q_{nm} = f_{nm} \Delta z \Delta v$. In order to track the evolution of the distribution function, each super-particle is tracked continuously in time (Lagrangian frame of reference). The value of the distribution function on the original grid points

can then be determined via interpolation after which the current and charge densities can be computed via numerical integration. More generally, the distribution function can be thought of as having the functional form

$$f(z, v, t) = \sum_{i=1}^{N_z N_v} w_i S_v(v - v_i(t)) S_z(z - z_i(t)) \quad (16)$$

where the summation index i is over all super-particles. The quantities S_v and S_z are shape functions and are an alternative means of expressing the interpolation process. The trajectory $(z_i(t), v_i(t))$ of the i -th particle is determined via the Lorentz force. The quantity w_i represents the particle weight that comes from the initial distribution function. The current density is then computed via

$$J(z, t) = \int v f(z, v, t) dv. \quad (17)$$

Another popular semi-Lagrangian scheme follows a procedure similar to Gibby et al. (2008) and Gibby (2008) in which the particles are only traced backwards for one time step and are not tracked continuously in time. This may result in some artificial diffusion, however, the number of super-particles that are in a cell at any given time are always known, which reduces the computational cost relative to the method of Nunn (1990).

Another class of methods which has not been utilized significantly for non-linear wave-particle interactions modeling are fully Eulerian schemes (Sonnendrücker et al., 1999; Harid et al., 2014b). In this approach, the Vlasov equation is solved numerically as a PDE using a finite difference or finite volume formalism. The advantage of such a technique is that the simulations are relatively simple to code and the stability criteria are well understood. Nevertheless, the grid based Courant-Friedrichs-Lewy (CFL) condition can be quite stringent and a non-uniform grid is difficult to implement. However, such methods can be successfully utilized with appropriate curvilinear coordinate transformations (Harid, 2015).

All the aforementioned VCON methods have only been considered for the whistler instability by using the narrowband approximation of Maxwell's equations (6,7). An important extension of this work for future researchers would be to utilize a fully broadband formalism for the wave equations along with a VCON solver.

4.2.2. Particle Methods

PIC codes, on the other hand, do not rely on a phase-space grid and continuously track the particle trajectories through time. Particle based techniques are thus often referred to as Lagrangian schemes. The currents are calculated by assuming a shape to the "super-particles" and accordingly interpolating to the spatial grid points where the wave fields need to be calculated. Mathematically, the distribution function in a PIC simulation can be written as

$$f(z, v, t) = \sum_{i=1}^{N_z N_v} w_i \delta(v - v_i(t)) S_z(z - z_i(t)) \quad (18)$$

The expression in (18) is essentially the same as (16), with $S_v(v - v_i) = \delta(v - v_i)$. That is, the shape function in velocity space is modeled as a Dirac delta function. This subtle feature allows for a significant reduction in computational resources. The computational burden is relieved when computing the current and charge densities since the delta function makes the velocity integrals trivial. Specifically, the current density is given by

$$J(z, t) = \int v f(z, v, t) dv = \sum_{i=1}^{N_z N_v} v_i w_i S_z(z - z_i(t)) \quad (19)$$

The lack of a velocity grid is a salient feature of PIC codes that is computationally desirable. However, although the computational cost of high-dimensional grid generation is removed, PIC simulations in turn suffer from numerical noise due to the random sampling of particles. For this reason, PIC codes often require millions of particles to reduce the artificial noise that is introduced (Birdsall and Langdon, 2004). Even so, current computational resources have permitted the use of PIC simulations to model the whistler mode instability. The works by Katoh and Omura (2008) and Hikishima and Omura (2012) demonstrate the clear utility of modern PIC simulations with a promising outlook for future computer experiments.

5. FUTURE WORK

5.1. Theory and Simulations

Over the past decades several numerical simulations have been utilized to clarify several aspects of wave-particle interactions in the magnetosphere. Even so, many observed phenomena have yet to be properly understood and certain physical assumptions used in current simulations may need to be relaxed. For instance, most models have primarily considered whistler mode interactions, yet electrostatic instabilities are believed to play an important role in the wave-particle interactions process (Omura et al., 2009). Particularly, the gap in chorus wave energy at frequencies around the half gyrofrequency may be in part due to Landau damping. The formation of an electron hole in phase-space inherently introduces non-zero space charge density that can drive quasi-electrostatic fields. Additionally saturation of whistler mode signals may also be due to mode conversion with electrostatic waves (Nunn, 1974). Effects known as wave-wave scattering involving interactions between electrostatic, electromagnetic, and quasi-electrostatic (lower hybrid) modes in the magnetosphere (Ganguli et al., 2010; Crabtree et al., 2012) are also yet to be explored with self consistent models.

An important simplification that is often used in simulation is ducted propagation. The waves are believed to be guided by field-aligned density irregularities that effectively force waves to propagate parallel to the geomagnetic field lines. However, additional spatial effects due to a finite sized guiding structure has not been explicitly considered in most modeling efforts. Since the ducts effectively act like a waveguide, they may also be responsible for exciting additional plasma modes via waveguide mode conversion. In the absence of ducts, higher dimensional effects may still be important as waves

propagate away from the equator (Ke et al., 2017). Thus, including two and eventually three dimensional features in space would serve as an important contribution to magnetospheric research.

The ideal simulation would consider a full six-dimensional model of the particle distribution in phase-space while self-consistently modeling the wave evolution in all three spatial dimensions. Current state-of-the-art computational resources may still be inadequate to solve the general problem. However, incrementally introducing additional physics will help isolate and clarify the dominant physical phenomena behind non-linear wave-particle interactions in the magnetosphere.

5.2. Experimental

Almost 30 years have past since the dismantling of Siple Station and there are currently no plans known to the authors to construct new facilities for radiation of ELF and low VLF waves with conventional multi-kilometer antennas. Future work on wave injection from the ground will therefore most likely take place at the HAARP facility. Although the management of the HAARP facility moved from the US Air Force to the University of Alaska in 2014, the facility continues to be used for active heating experiments. New formats can be designed and transmitted to validate numerical simulations that have greatly increased their capabilities in the last few years. Active experiments can serve not only to shed light on the fundamental theoretical process of non-linear cyclotron resonance and its dynamic evolution but also provide practical results for future strategic schemes of radiation belt mitigation (Inan et al., 2003). In the latter the key objective is pitch angle scattering, which strongly depends on wave amplitude. Therefore learning under what conditions what frequency-time formats can lead to the greatest amplification is a key question. Such investigations would build upon the identification of positive frequency-time ramps as being favorably amplified in a large number of past experiments. Another outstanding question in magnetospheric physics that active experiments are ripe to address is the relationship between natural hiss and chorus waves. In the past it has been proposed that either hiss creates chorus (Koons, 1981) or chorus is the source of hiss (Bortnik et al., 2008). Experiments can be performed to test whether hiss like signals can trigger chorus or how spectral changes of coherent injected signals can evolve to appear like hiss emissions. The already completed preliminary investigations of coherence bandwidth done at HAARP (Golkowski et al., 2011) show that this is a fruitful line of investigation. The real-time interaction of the injected signals with present chorus and hiss emissions is also worthy of deeper investigation.

Wave injection and subsequent observation of whistler mode waves from the ground requires at least minimal guiding along the geomagnetic field from density irregularities or the plasmapause. The presence or absence of these structures has been shown to affect the occurrence of observations (Golkowski et al., 2011). There have been efforts to use the HAARP facility to generate field aligned irregularities that could guide waves to the conjugate point (Milikh et al., 2008). At the same time,

whether or not such efforts create structures that extend along the entire field line and can compete with processes in the natural environment has been called into question by other authors (Pidtychiy et al., 2011). More investigations in this area seem appropriate. For all of the studies proposed there is no question that a conjugate observation station for HAARP with a ELF/VLF receiver and other instruments would be extremely useful for wave injection experiments and also other investigations performed at HAARP.

One aspect of past studies that has only seen mixed results in the Siple and HAARP experiments is the detection of transmitter induced energetic electron precipitation from the magnetosphere. As discussed above, the interaction leading to amplification is a manifestation of the same fundamental process as pitch angle scattering. However, detection of energetic electron precipitation from the magnetosphere on the ground is challenging and often involves indirect methods. Direct one to one correlation between precipitation signatures and Siple Station transmissions has not been reported even though X-ray observations on balloon platforms have shown evidence of precipitation from individual chorus elements (Rosenberg et al., 1981). The recent work of the Balloon Array for Radiation belt Relativistic Electron Losses (BARREL) mission (Woodger et al., 2015) and FIREBIRD II cubesat (Breneman et al., 2017) have shown that balloon platforms and small satellites can be effective tools in observing energetic electron precipitation going forward. At HAARP an attempt was made to detect induced precipitation using VLF remote sensing and also the Poker Flat incoherent scatter radar (ISR) but did not lead to conclusive findings in the limited attempts that were made (Golkowski, 2009). Increasing ISR capability at the HAARP facility is seen as the best way to approach future induced precipitation studies. It is noted that evidence of direct precipitation induced from a VLF ground transmitter was reported in the SEEP experiment (Imhof et al., 1983; Inan et al., 1984). At the same time, more recent experiments with the NPM transmitter in Hawaii although initially interpreted as bearing evidence of precipitation were later shown to be more ambiguous (Graf et al., 2011). A thorough investigation of controlled precipitation using multiple detection methods could have a broad impact on numerous efforts in the magnetospheric community.

The upcoming US Air Force *Demonstration and Science Experiments* (DSX) mission brings with it the exciting prospect of controlled radiation of waves directly in the magnetosphere (Scherbarth et al., 2009). Space based injection will have easier access to the non-linear wave-particle regime since the high losses from penetration of the ionosphere will be absent. Observationally, full disambiguation of the non-linear growth process would require multiple closely spaced satellites to observe transmissions along their propagation path. In this context closely spaced (< 200 km) spacecraft observations have been shown to be very fruitful for investigations of chorus wave properties (Santolik and Gurnett, 2003). Deploying such spacecraft for wave injection observations would be most effective if the spacecraft could be arranged to be along the same geomagnetic field line.

6. CONCLUSION

Controlled excitation of non-linear whistler mode wave particle interactions has a rich and fruitful history. Experimental activities and likewise the theoretical and computational efforts they have motivated have been a cornerstone of near-Earth space physics. The current times embody quickly improving computational tools and ever easier access to space with improved sensors and hardware capabilities. Conditions are thus favorable for active controlled experiments to yield new fundamental discoveries.

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The Effect of Plasma Releases on Equatorial Spread F —a Simulation Study

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A currently unfulfilled goal of active experimentalists is to control the occurrence of instabilities in the ionosphere such as Equatorial Spread F (ESF), which generates large-scale electron density depletions (plasma bubbles) in the night-time ionosphere at low latitudes. It has been theorized that by artificially injecting ionizing chemicals (such as barium) into the ionosphere, ESF may be suppressed. Large plasma releases modify the ionospheric conductance, which affects the electrodynamics of the system and may thereby influence the growth (or suppression) of ESF. In this study, the feasibility of controlling ESF growth via plasma releases in the ionosphere is examined for the first time using a fully three-dimensional, first-principles model of the ionosphere: SAMI3/ESF (Sami is Another Model of the Ionosphere). The numerical simulations show that under certain circumstances plasma injections may be able to trigger or suppress ESF growth. The results indicate that the plasma density must be above a threshold level to sufficiently modify the ionospheric conductance. In addition, the plasma must be injected at a suitable location and time. The results of this numerical investigation provide guidance for future experimental campaigns.

Keywords: ionosphere, equatorial spread F , active experiment, chemical release, equatorial plasma bubble

1. INTRODUCTION

The phenomena of Equatorial Spread F (ESF) has long been of interest to the aeronomy community (e.g., Farley et al., 1970; Ossakow, 1981; Hysell, 2000; Woodman, 2009; Abdu, 2012) because these large-scale perturbations have serious equatorial space weather implications such as disruption in radio communication, navigation, and geo-positioning. They also influence the performance and reliability of space borne and ground based electronic systems. It commonly occurs in the post-sunset ionosphere when E region conductivity drops and the equatorial F region ionosphere can become unstable because of a Rayleigh-Taylor (R-T) like instability (e.g., Sultan, 1996). These internally driven perturbations occur naturally on a day-to-day basis in the low-latitude region. These instabilities can generate large scale (10 Km), low density (1–3 order of magnitude smaller) plasma bubbles that can ascend to 1,000 Km at their apex. Such equatorial depleted plasma density regions are of great interest to the space weather community because they can also extend into the middle latitudes along the geomagnetic field lines (e.g., Huang et al., 2007; de La Beaujardière et al., 2009) and interfere with the operation of space borne and ground based technological systems.

First detection of an ESF event was reported in Berkner and Wells (1934) using an ionosonde. Woodman and La Hoz (1976) reported the first plasma depletion (plasma bubbles) detection. Since then ESF and associated bubbles have been extensively studied with an armada of ground-based (e.g., Tsunoda, 1983; Kudeki and Bhattacharyya, 1999), space-based (e.g., Burke et al., 2003; Le et al., 2003; Kelley et al., 2009; Huang et al., 2011), *in-situ* studies with CHAMP (Stolle et al., 2006) and SWARM (Wan et al., 2018), rocket measurements (e.g., Kelley et al., 1986; LaBelle and Kelley, 1986; Caton et al., 2017), and modeling studies (e.g., Zalesak and Ossakow, 1980; Zalesak et al., 1982; Huba et al., 2009a,b; Su et al., 2009; Krall et al., 2010a,b; Retterer, 2010a,b). Despite the decades of intensive research that have dramatically increased our current understanding of ESF bubbles, much about their triggering, suppression, and capricious day-to-day variability in occurrence is poorly understood. Plasma density perturbations (“seed perturbations”) generated through non-linear hydrodynamical R-T instability in the bottomside *F* region ionosphere are considered the cause of ESF bubbles, and thus are one of the primary seeds used in modeling studies (e.g., Retterer and Roddy, 2014). The seed perturbations for plasma bubbles can be associated with atmospheric gravity waves (e.g., Huang et al., 1993; Kudeki et al., 2007; Abdu et al., 2009; Takahashi et al., 2009). Another mechanism of seed perturbation in the bottomside *F* region is associated with vertical shear in zonal plasma drifts (e.g., Kudeki and Bhattacharyya, 1999; Hysell et al., 2005, 2006; Kudeki et al., 2007). The ESF bubble development is directly dependent on the magnitude of initial perturbation and magnitude of R-T growth rate (e.g., Retterer and Roddy, 2014). The evolution of ESF bubbles is complex; therefore, there is a common consensus that numerical simulation experiments in addition to observational experiments are necessary to understand their formation and evolution.

The present study is focused on a numerical plasma seeding investigation of controlling ESF growth (triggering and suppression) by inserting plasma (such as in artificial ionospheric modification experiments discussed in e.g., Bernhardt, 1992; Huba et al., 1992; Caton et al., 2017; Retterer et al., 2017) along a magnetic field line at different altitudes that modifies ionospheric conductance. For this numerically controlled investigation, we utilize the capabilities of SAMI3/ESF model that has been used in a number of other studies to investigate ESF (e.g., Huba et al., 2009a; Krall et al., 2009; Zawdie et al., 2013). To our knowledge, this is the first numerical diagnosis to control the ESF growth self-consistently from first principles. This study is motivated by the results from earlier controlled ionospheric modification experiments for tailoring radio wave propagation medium (such as Wright, 1964; Pickett et al., 1985; Çakir et al., 1992; Caton et al., 2017; Retterer et al., 2017) by perturbing ionospheric densities through chemical or plasma releases. The results of this numerical study enhances our current understanding that is required for ionospheric modification efforts to control the ESF bubbles using plasma releases.

2. PREVIOUS WORK

The idea of artificially inducing equatorial spread *F* has a long history. Ossakow et al. (1978) performed the first set of simulations to demonstrate that a large plasma depletion in the bottomside, equatorial ionosphere after sunset could trigger large scale plasma bubbles, similar to those observed during naturally occurring equatorial spread *F*. Subsequently, The Brazilian Ionospheric Modification Experiments (BIME) was carried out in September 1982 (Klobuchar and Abdu, 1989). Nike-Black Brant rockets were launched from Natal, Brazil on separate evenings. The rockets injected H₂O and CO₂ into the bottomside *F* layer to create an artificial electron density depletion. This was done after sunset when the ionosphere was rising. These depletions were tracked moving eastward using TEC (Total Electron Content) measurements and oblique ionosondes. Subsequently, plasma bubble irregularities were detected as spread *F* echoes, scintillation of radio waves, and TEC depletions roughly 300 – 500 km east of the injection site. Equatorial spread *F* was not observed on other nights of the campaign suggesting that plasma bubbles were generated artificially by the chemical releases.

Another chemical release experiment was carried out as part of the NASA Combined Release and Radiation Effects Satellite (CRRES) mission to induce equatorial spread *F* (Sultan and Jared, 1994). Sulfur hexafluoride (SF₆) was released into the bottomside *F* layer from sounding rockets launched from Kwajalein. Several diagnostics were used to monitor the ionosphere before and after the releases [e.g., incoherent scatter radar, High Frequency (HF) radar, and optics]. Small equatorial spread *F* plumes were observed during the experiments suggesting that they were artificially induced.

A recent chemical release experiment was launched in May of 2013 as part of the Metal Oxide Space Cloud (MOSC) experiment by the Air Force (Caton et al., 2017). In this experiment clouds of vaporized samarium were released from sounding rockets launched from the Reagan Test Site in Kwajalein Atoll. A numerical experiment examined the electrodynamic effects of the plasma clouds produced by the MOSC campaign (Retterer et al., 2017). The study was able to reproduce a “comma-like” flow around the cloud that was observed during the experiment. The simulations also suggested that if the MOSC plasma clouds were denser and closer to the bottom edge of the *F* region, they may have been able to suppress the development of ESF.

A related study has also been performed using the SAMI3/ESF model, which investigated whether ESF bubbles could be triggered with artificial HF (High Frequency) radio wave heating (Zawdie and Huba, 2014); it demonstrated that the density perturbations due to artificial HF heating of the ionosphere would not generate ESF bubbles. The artificial HF heating increases the electron temperature causing a pressure gradient that drives electrons down the field lines to higher latitudes away from the heating location. Since the artificial HF heating redistributed electron density along the field line, the Pedersen conductance and ESF growth rate were not significantly affected by the HF density perturbation and ESF was not triggered. In order to trigger ESF, the field line integrated electron density

would need to be modified, which does not occur during HF heating.

This numerical experiment is targeted at determining the feasibility of both triggering and suppressing ESF bubbles in the ionosphere via plasma releases in the ionosphere. This is the first time a fully 3D first-principles ionospheric model (SAMI3/ESF) has been employed for investigation of ESF control using artificial plasma injections.

3. MODEL DESCRIPTION

In the present study, we use SAMI3/ESF; a full description of the model can be found in Huba et al. (2008, 2009b). Here, we highlight the main features and modifications of the model used. SAMI3/ESF is a three dimensional physics-based ionosphere model based on SAMI2 (Huba et al., 2000). SAMI3 simulates the ionospheric plasma on a nonorthogonal, nonuniform grid that follows the dipole electric field lines. It can be run in a global mode, or in a wedge mode for high resolution studies of ESF. In wedge mode, the SAMI3 grid consists of a narrow wedge in longitude (about 4°). SAMI3/ESF includes seven ion species (H^+ , He^+ , N^+ , O^+ , N_2^+ , NO^+ , O^+). The ion continuity and momentum equations are solved for all seven species; in addition, the temperature equations are solved for H^+ , He^+ , O^+ , and the electrons. Quasi-neutrality is assumed, so the electron density is simply the sum of densities of the ion species. It self-consistently calculates the electric potential, which is used to calculate the $E \times B$ drifts in the perpendicular (vertical and longitudinal) directions. For simplicity, in the present case study, we use non-tilted dipole magnetic field, which means magnetic and geographic latitude are the same.

To simulate the effects of a plasma release on the ionosphere, we have added an eighth species (Al^+) to the model. The ion continuity, momentum and temperature equations are solved for the additional ion. Since aluminum is nonreactive, it is assumed that there are no significant chemical reactions with the other seven ions. The addition of the new ions, however, does significantly affect the ionospheric conductivities and electrodynamics, as will be examined in the following section. The aluminum is assumed to be ionized at the time of release, and the initial release is a Gaussian distribution:

$$n_{Al^+} = (n_0) e^{[-(s-s_0)^2/\Delta s^2]} e^{[-(p-p_0)^2/\Delta p^2]} e^{[-(\phi-\phi_0)^2/\Delta \phi^2]} \quad (1)$$

where, s is the direction along the field line, p is in the direction of increasing field lines, and ϕ is in the longitudinal direction. The Gaussian is initialized in the dipole coordinates. For a nominal simulation, n_0 is the initial release density of $5.0 \times 10^7 \text{ cm}^{-3}$, Δs is 15 km, Δp is 7 km and $\Delta \phi$ is 7 km. The parameters s_0 , p_0 , and ϕ_0 define the release location, which varies depending on the simulation case. In each simulation, the aluminum initial release is added a few time steps into the simulation, the ions subsequently evolve self-consistently according to ion continuity, momentum and temperature equations, which allows their effect on the growth of ESF bubbles to be examined.

For initialization, SAMI3/ESF uses output from a 48 h run of the SAMI2 model. For the initial conditions, SAMI2 was

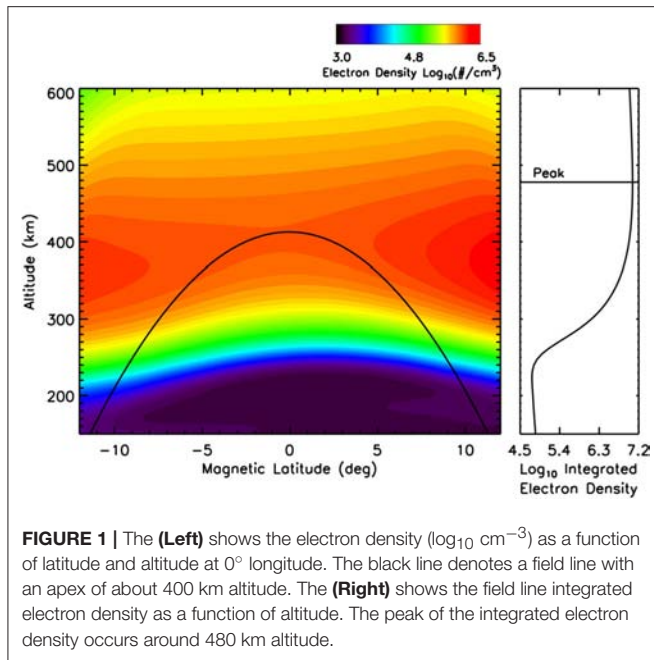
run using the following conditions: F10.7 = 100, F10.7A = 100 (F10.7A is the 81-day centered average of F10.7), $A_p = 4$, and day of year = 130. The plasma parameters at 19:30 Universal Time (UT) of the second day were used to initialize the three dimensional model. The initialization parameters are consistent with earlier studies that used the SAMI3 model (Huba and Krall, 2013; Zawdie and Huba, 2014) in order to ensure that the background conditions are sufficient for ESF generation. Previous studies have shown that ESF bubbles simulated with SAMI3 match well with observations (e.g., Krall et al., 2009, 2010b); these comparisons are not examined in detail in this work. In addition, while the simulation parameters are consistent with observations (Stolle et al., 2006; Yizengaw and Groves, 2018), the daily, seasonal, and longitudinal variability of ESF are not considered in this paper. Because the purpose of study is to understand the behavior of local plasma features by ionospheric modification, SAMI3/ESF is run in the wedge mode rather than global. Due of the local nature of the study, we have not included the effect of thermospheric winds on the results.

4. SIMULATION RESULTS

4.1. Plasma Releases: The Basics

ESF bubbles are a Rayleigh-Taylor like instability, where a dense fluid lies on top of a lighter fluid and small perturbations become unstable. In the ionosphere, this occurs after dusk when the F1 and E-regions rapidly recombine, leaving a heavier F2 region at higher altitudes. Due to the ionospheric electrodynamics, this instability occurs along full magnetic field lines, so the localized electron density is less important than the total electron density integrated along the magnetic field line. The daily, seasonal and longitudinal variability of ESF bubble occurrence are not yet fully understood. Recent work has demonstrated that a wide variety of geophysical parameters may be important in predicting the timing and locations of ESF development (e.g., Stolle et al., 2006; Carter et al., 2014; Retterer and Roddy, 2014; Yizengaw and Groves, 2018). The effect of geophysical parameters on the development of ESF are not investigated in this work; instead, this study investigates how changes to the plasma density affect the ionospheric electrodynamics.

First, we performed a simulation with the SAMI3/ESF model without any perturbations as a background case. Then, a number of simulations were performed with the SAMI3/ESF model with simulated plasma releases at different locations in the ionosphere in order to determine their effect on the creation/inhibition of ESF bubbles. The electron density as a function of latitude and altitude at 0° longitude for the background simulation is shown in the left panel of **Figure 1**. The right panel of **Figure 1** shows the field-line integrated electron density for the background simulation as a function of the field-line apex altitude. Note that the peak electron density at 0° latitude occurs around 400 km altitude, but the peak of the integrated electron density is around 480 km and is marked in the right panel of **Figure 1**. The general approach to simulate ESF bubbles in physics-based models is to add a small perturbation, or seed to a field line slightly below the peak field-line integrated electron density. This seed triggers the instability, growing into a bubble extending along the field line

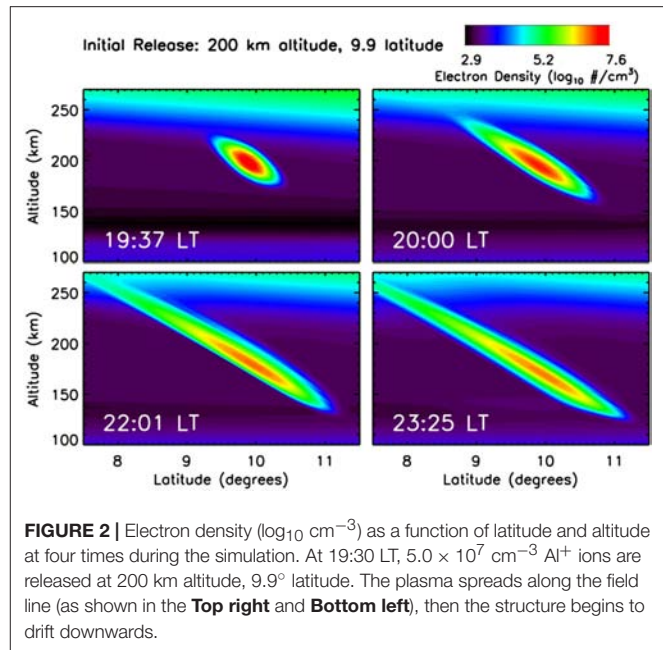


and lifting up through the *F* region ionosphere. **Figure 1** (left panel) shows such a field line outlined in black. For this study, plasma releases were also added to this field line to determine their effect on ESF bubble development.

Figure 2 shows an example of the evolution of a plasma release in the ionosphere. At 19:36 local time (LT) a Gaussian blob of $5.0 \times 10^7 \text{ cm}^{-3}$ Al^+ ions are released into the simulation at 200 km altitude, at location 9.9° latitude and 0° longitude. Over the next two and a half hours, the ions spread along the field line, extending between 150–300 km altitude. In addition, the extended cloud drifts downwards. **Figure 3** shows a similar example of a plasma release, but at the apex of the field line (400 km altitude, 0° latitude, 0° longitude). The ions quickly fall down along the field line, extending out to $\pm 2^\circ$ latitude within 30 minutes. The cloud also begins to drift downward due to the polarization electric field generated by the plasma cloud. Generally, the larger the plasma release is, the stronger the polarization field becomes, so the denser a plasma cloud is the more quickly it will fall. It should be noted that although these particular simulations used Al^+ as the ion species, similar simulations have been performed with Lithium ions and the results were qualitatively the same.

Figures 2, 3 show releases of Aluminum ions at different locations along the same field line. Although the growth rate of an ESF bubble is dependent on field line integrated quantities, the altitude where the plasma release occurs drastically affects both the Pedersen conductance and growth rate. The growth rate of an ESF bubble can be calculated as in Zawdie and Huba (2014):

$$\gamma = -\frac{\int \sigma_{Hc} (g_p/L_n) ds}{\int \sigma_p ds} \quad (2)$$



where g_p is the gravitation term, s is in the direction of the field line, $L_n^{-1} = \partial \ln(n_0)/\partial p$, n_0 is the electron density, p is perpendicular to the field line, σ_p is Pedersen conductivity, and σ_{Hc} is Hall conductivity. The Pedersen and Hall conductivities can be approximated as:

$$\sigma_p \approx \sum_i \frac{nec}{B} \frac{v_{in}}{\Omega_i} \quad (3)$$

$$\sigma_{Hc} \approx \sum_i \frac{nec}{B} \frac{1}{\Omega_i} \quad (4)$$

where v_{in} is the ion-neutral collision frequency, $\Omega_i = eB/m_i c$, n is the electron density, e is the electron charge, c is the speed of light, B is the magnetic field strength, and m_i is the mass of ion i . In Equation 2, $\int \sigma_p ds$ is the Pedersen conductance; thus the growth rate is inversely proportional to the Pedersen conductance.

Figure 4 shows the Pedersen conductance and maximum growth rate as a function of the release altitude of the plasma injection along the field line outlined in **Figure 1**. The dotted line denotes the Pedersen conductance and growth rate for the background simulation case (no plasma release). The top panel shows that the lower the release altitude, the larger the increase in Pedersen conductance. Plasma releases above 300 km altitude do not have a substantial effect on the Pedersen conductivity. On the other hand, the bottom panel shows that the higher the release altitude, the larger the increase in the ESF growth rate. Plasma releases below 250 km do not significantly increase the ESF growth rate. Based on the Pedersen conductivity and growth rate change with altitude, we selected our plasma release altitudes along the selected magnetic field line for ESF bubble control. The following sections describe the plasma

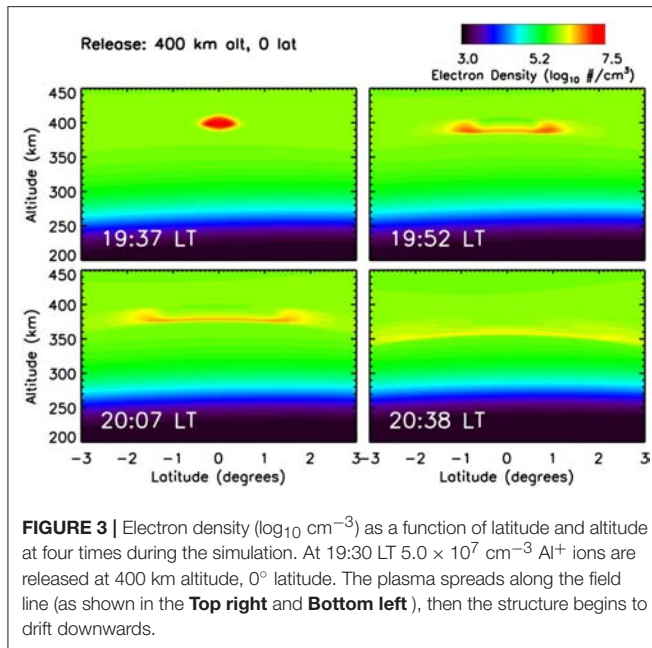


FIGURE 3 | Electron density ($\log_{10} \text{cm}^{-3}$) as a function of latitude and altitude at four times during the simulation. At 19:30 LT $5.0 \times 10^7 \text{ Al}^+$ ions are released at 400 km altitude, 0° latitude. The plasma spreads along the field line (as shown in the **Top right** and **Bottom left**), then the structure begins to drift downwards.

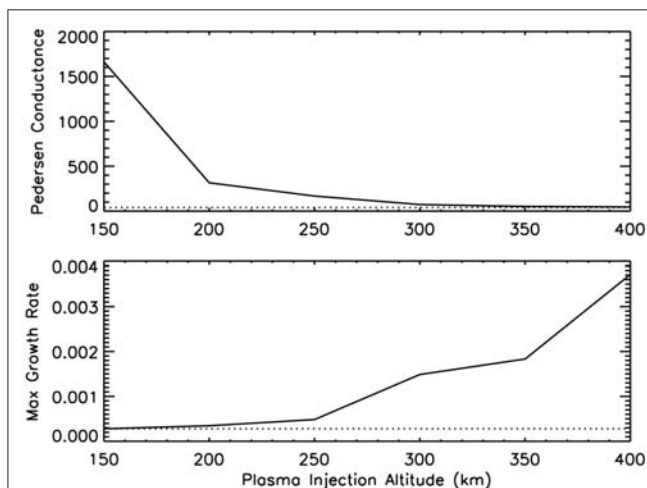


FIGURE 4 | The **Top** shows the Pedersen conductance (mho) of the field line outlined in the left panel of **Figure 1** for simulations with plasma releases at different altitudes. In each simulation $5.0 \times 10^7 \text{ Al}^+$ ions were released at a different location along the field line. The dotted line denotes the Pedersen conductance for a simulation with no plasma release. The **Bottom** shows the maximum growth rate (s^{-1}) for a the same series of simulations as the **Top**. The dotted line shows the maximum growth rate for a simulation with no plasma release.

release simulations designed to test whether the perturbation of the Pedersen conductance/growth rate can suppress/trigger ESF bubbles.

4.2. How to Trigger ESF Bubbles

As shown in **Figure 4**, a plasma release at the apex of the field line significantly increases the ESF growth rate. In this section, the results of a simulation where plasma is released

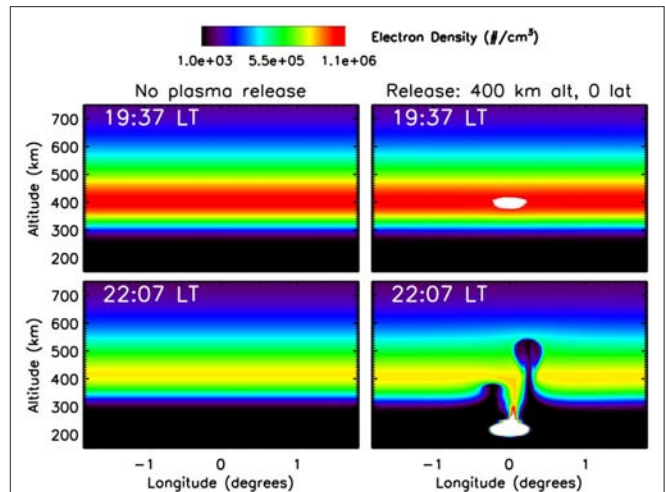
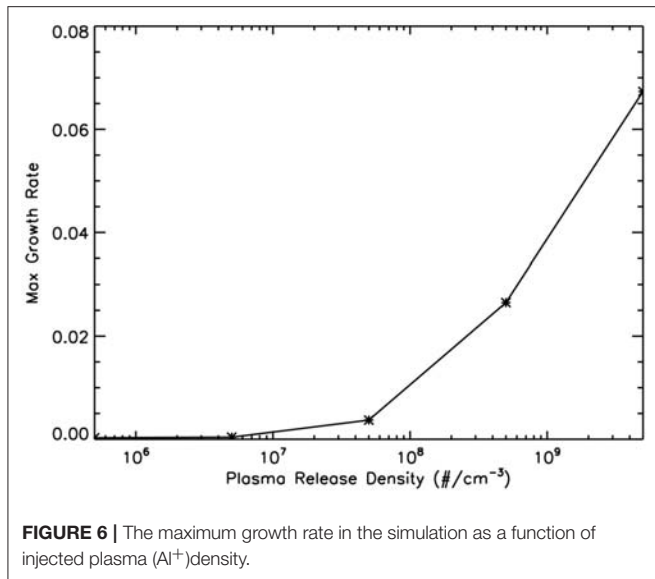


FIGURE 5 | Electron density ($\log_{10} \text{cm}^{-3}$) as a function of longitude and altitude at 0° latitude for two different simulations (**Left, Right**) at two different times (**top and bottom**). The left panels show the time evolution for the background simulation with no plasma release; in this case no ESF develops. The right panels show the case where a plasma release is simulated at 400 km altitude and 0° latitude. Note that the color bar saturates in the area of the plasma release. As shown in **Figure 3** the plasma release spreads along the field line and then drifts downward. The bottom right panel shows that after several hours ESF bubbles develop in the ionosphere as a result of the plasma release.

at the apex ($\sim 400 \text{ km}$) of the seeding magnetic field line, are examined to determine if the increase in growth rate is sufficient to trigger the growth of an ESF bubble. The results of this numerical case study are shown in **Figure 5**, which shows the electron density as a function of longitude and altitude at 0° latitude. The left panels show the ambient electron density from the background simulation with no plasma release. The right panels show the simulation results for a release of $5.0 \times 10^9 \text{ cm}^{-3} \text{ Al}^+$ at 400 km altitude. In the top right panel, the plasma release is seen as a white blob (the color scale is saturated) at 400 km altitude just after the release. The lower panels show the evolution of the system after two and a half hours. In the background case (left), no ESF bubble forms, but in the plasma release simulation (right), an ESF bubble has been triggered and is rising through the ionosphere. Thus, by releasing plasma at the apex of the field line, an ESF bubble has been triggered.

Further investigations have indicated that there is a minimum plasma release density required to trigger an ESF bubble. **Figure 6** shows the maximum growth rate that results from plasma releases of different densities at 400 km altitude. The larger the plasma release density, the larger the maximum growth rate. A plasma release of $5.0 \times 10^7 \text{ cm}^{-3} \text{ Al}^+$ was found to be sufficient to generate ESF, but the bubble formed more slowly in the ionosphere (3.5 h) than with a release of $5.0 \times 10^9 \text{ cm}^{-3} \text{ Al}^+$ (2.5 h). In our investigation, we also found that a plasma release at a lower altitude (such as 350 km) could trigger ESF, but the density threshold increases as the release altitude decreases. The numerical case study of the



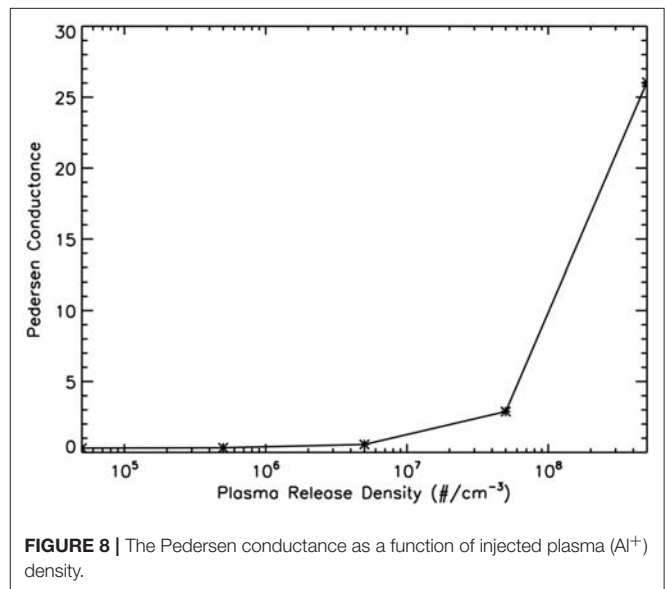
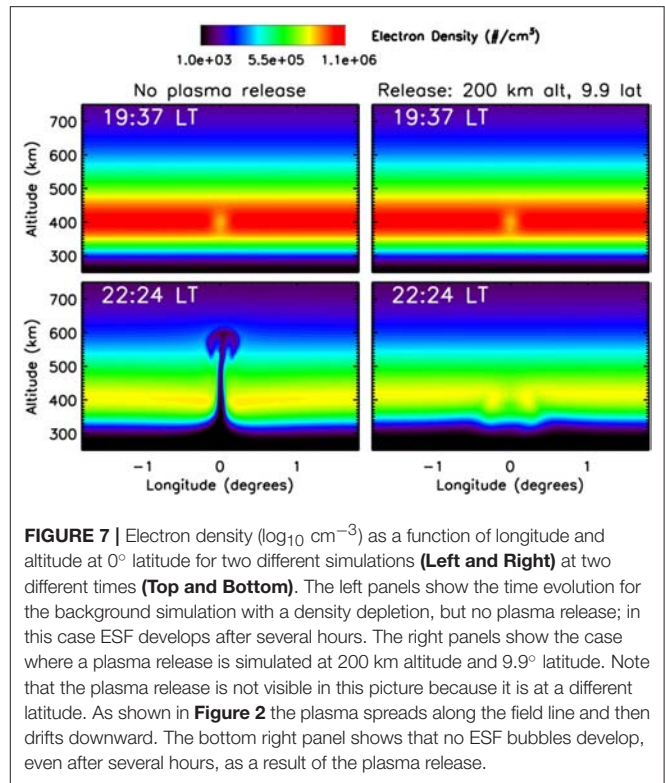
plasma releases below 250 km altitude are examined in the following section.

4.3. How to Suppress ESF Bubbles

Figure 4 shows that adding a plasma release at a lower altitudes along the seeding field line can substantially increase the Pedersen conductance. The ESF growth rate is inversely proportional to the Pedersen conductance, so it is not unreasonable to suggest that a plasma release at a low altitude (around 200 km) may suppress ESF. In this section, simulation results are examined to determine if it is feasible to suppress ESF bubbles via plasma releases. Two simulations were performed; the first had a density depletion along the seeding field line that is typically used to trigger ESF bubbles in simulations. The second had a density depletion and a plasma release of $5.0 \times 10^7 \text{ cm}^{-3}$ Al^+ ions at 200 km altitude, 9.9° latitude, and 0° longitude. **Figure 7** shows the results of the two simulations: the electron density as a function of longitude and altitude. The left panels show the results with no plasma release and the right panels show the result of injecting plasma.

The top two panels show the results shortly after the simulations begin. In both the ambient and plasma release cases, an electron density depletion in the seeding field line at 400 km altitude, 0° longitude is clearly visible. Note that the plasma release does not appear on this plot because it is at a different latitude (9.9° latitude) not covered in this figure. The bottom two panels show the results after 3 h. The simulation case with only the density depletion (left) shows a well developed ESF bubble. The case with the density depletion and the plasma release shows some density perturbations in F region around 0° longitude, but the ESF bubble has successfully been suppressed.

Further simulations indicate that there is a minimum size and density for a plasma release to suppress such a ESF bubble. In order to suppress an ESF bubble, the plasma release must cover the field lines associated with the density depletion that seeds



the bubble. In this case, the density depletion and the plasma release both covered 15 and 35 km in the longitude and latitude directions, respectively. The larger the density depletion is, the larger the plasma release needs to be. The plasma release also needs to be in a suitable location in order to affect all field lines that have been perturbed by the density depletion. In addition to these constraints, there is a threshold (lower limit) where the plasma release is not dense enough to suppress an ESF bubble.

This is depicted in **Figure 8**; it shows the Pedersen conductance as a function of plasma release density, which increases with decreasing plasma release altitude. A release of less than $5.0 \times 10^6 \text{ cm}^{-3} \text{ Al}^+$ ions does not increase the Pedersen conductance enough to suppress the ESF bubble. It should be noted that this threshold is lower for releases at lower altitudes, although a release below 150 km has not been examined by this study.

5. CONCLUSIONS

The numerical investigations of the present study suggest that plasma releases can be used to trigger and suppress ESF bubbles. These results are achieved by injecting plasma at different locations along a key “seeding field line” which has been used in simulations to trigger ESF. This particular field line is located below the peak of the field-line integrated electron density. Plasma releases at or near the apex of the magnetic field line spread down the field line, then the cloud drifts down to lower altitude; in the process, an ESF bubble is triggered. A plasma release at a lower altitude along the same field line can suppress an ESF bubble, provided that the plasma cloud covers all field lines affected by the ESF seed. The key challenges to manipulating the growth of ESF are ensuring that the initial plasma release is dense enough to affect the Pedersen conductance or growth rate and that the release occurs in the correct place. One obvious question is: how practical are these mechanisms for controlling ESF bubbles?

Based on the results of this numerical investigation, triggering ESF bubbles via high altitude plasma releases is likely a feasible project, however, the key engineering issue is ensuring that the injected plasma density is larger than $5.0 \times 10^9 \text{ cm}^{-3}$ ions. In order to trigger an ESF with plasma injection, the plasma must be injected at suitable location (magnetic equator) and time (19:30 LT) as we found in our test cases. Since the plasma cloud creates a polarization electric field and drifts down in altitude, the only difficulty would be ensuring that the release occurs above the F-peak and that it is timed correctly. If the release occurs at too low of an altitude, it may fail to trigger the ESF, but as long as the release occurs above the F-peak, the cloud should drift down through the region where it can trigger ESF.

Suppressing ESF in the ionosphere is likely more difficult. In addition to the constraint that the plasma releases be dense enough to adequately modify the Pedersen conductance, there are also significant limitations on the location of the plasma injection. In particular, the injection must create a plasma cloud that extends over all field lines affected by potential ESF triggers. In our numerical experiment the plasma cloud was 15 km in the longitudinal direction by 35 km in the latitudinal direction,

but that only worked because the exact location and size of the density depletion was known. It is possible that with additional measurements, one could determine the most likely position of a seeding field line in a particular longitude sector. Then it is necessary that the plasma cloud extend in latitude and longitude enough to cover any potentially affected field lines.

Another complication is the presence of the neutral wind in the thermosphere. The neutral wind directly affects the growth of ESF bubbles in the ionosphere as shown in Krall et al. (2009) and Huba and Krall (2013), but the neutral wind also affects the distribution of plasma releases in the ionosphere. This is primarily an issue for attempting to suppress ESF bubbles, as the plasma release may be driven by the neutral winds to other longitude sectors, making the determination of where to put the plasma release even more difficult. A full examination of the effect of neutral winds on plasma releases in the ionosphere and their effect on the growth/suppression of ESF is left for future work.

Our analysis is also relevant to extending and controlling the transverse to the magnetic field size of high kinetic beta ($\beta = 10^3 - 10^7$) plasma injections in the ionosphere using the capabilities of the ENIG Magneto-Hydrodynamic Flux Compression Generator (FCG) (Kim and Bentz, 2015). Initial tests of this new technology have been promising and demonstrate that it may be possible to control the size of a plasma release in the near future, potentially enabling technologies to suppress ESF bubbles.

DATA AVAILABILITY

The datasets generated and analyzed for this study available on request from the lead author.

AUTHOR CONTRIBUTIONS

KZ performed the model runs, generated the figures, and wrote the paper. JH assisted with the model simulations and wrote part of the paper. MD wrote part of the paper. KP guided the initial work and wrote part of the paper.

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Atmospheric Effects of a Relativistic Electron Beam Injected From Above: Chemistry, Electrodynamics, and Radio Scattering

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We present numerical simulations and analysis of atmospheric effects of a beam of 1 MeV electrons precipitating in the upper atmosphere from above. Beam parameters of 100 J or 1 kJ injected in 100 ms or 1 s were chosen to reflect the current design requirements for a realistic mission. We calculate ionization signatures and optical emissions in the atmosphere, and estimate the detectability of optical signatures using photometers and cameras on the ground. Results show that both instruments should be able to detect the beam spot. Chemical simulations show that the production of odd nitrogen and odd hydrogen are minimal. We use electrostatic field simulations to show that the beam-induced electron density column can enhance thunderstorm electric fields at high altitudes enough to facilitate sprite triggering. Finally, we calculate signatures that would be observed by incoherent scatter radar (ISR) and subionospheric VLF remote sensing techniques, although the latter is hindered by the limitations of 2D simulations.

Keywords: radiation belts, atmosphere, electron beam, chemistry, sprites, subionospheric VLF

1. INTRODUCTION

Electron guns firing artificial beams of electrons with energies in the tens of keV have been used since the 1970s to probe magnetospheric and auroral physics (e.g., Winckler, 1980; Neupert et al., 1982; Burch et al., 1993; Stone and Bonifazi, 1998). In the late 1980s, the idea of using a relativistic beam of electrons took hold (e.g., Banks et al., 1987, 1990); a comprehensive review of research into relativistic beam experiments up to 1992 was provided by Neubert and Banks (1992). Relativistic (MeV) beams have a variety of advantages over keV beams, including beam stability and reduced spacecraft charging effects (Neubert and Gilchrist, 2002), and faster propagation along magnetic field lines (Delzanno et al., 2016; Dors et al., 2017). Krause (1998) performed detailed calculations of the atmospheric response to a relativistic electron beam, including calculations of beam dynamics and stability, and ionization and X-ray production in the atmosphere. Further calculations were made by Neubert et al. (1996) on the propagation of the electron beam in the magnetosphere and the atmospheric response.

In recent years we have begun exploring a number of potential applications of space-based, artificial relativistic beam injection, including magnetic field line mapping, studies of wave-particle interactions, and studies of the atmospheric effects of precipitation, that should be enabled by the present state of technology in particle accelerators. However, any such experiment relies on

the ability to detect and measure the beam, either *in-situ* by directly observing the electron beam, or by remote sensing from the ground, by observing the atmospheric signatures of the beam precipitating in the atmosphere. For the latter, we need to assess the diagnostic signatures of the beam in the atmosphere.

Marshall et al. (2014) expanded on the work of Krause (1998) to calculate optical emissions observable from the ground, X-ray production and propagation and detectability from satellites and balloons, and backscattered electrons that could be observed from Low Earth Orbit (LEO). That study showed that optical signatures are likely detectable; indeed, the SEPAC experiments (Neubert et al., 1995) observed optical emissions of $\sim 1\text{--}5$ kR in the $4,278\text{ \AA}$ emission from an 1.2 A injected beam of 6.25 keV electrons, about a factor of 7.5 higher energy flux than our proposed 1 MeV beam. X-ray fluxes were likely to be far too low to be detectable from either LEO or balloon altitudes; and ionization could likely be measured from the ground using incoherent scatter radar. However, that study investigated a pulse of electrons with only $0.05\text{--}1$ Joules of total energy. Recent accelerator design efforts are targeting a beam total energy of $100\text{--}1,000\text{ J}$, prompting a revisit to the calculations of Marshall et al. (2014).

In this paper, we expand on the Marshall et al. (2014) study by increasing our simulated beam energy, and by investigating further atmospheric effects and diagnostic signatures of the relativistic electron beam injection. In particular, here we update our optical emissions and ionization calculations for a specific set of beam parameters; we calculate the chemical response of the atmosphere in terms of odd nitrogen and odd hydrogen production; we study the electrodynamic response of the atmosphere in the presence of thunderstorm electric fields; and we investigate subionospheric VLF remote sensing as a potential diagnostic of the electron density disturbance in the atmosphere. Together with Marshall et al. (2014), these results form a complete picture of the atmospheric response to an artificial beam injection, and calculate the expected response in all possible diagnostic methods.

2. MODEL INPUT PARAMETERS

In Marshall et al. (2014), we considered the beam-atmosphere interaction over a range of beam energies, but primarily focused on an electron energy of 5 MeV . In this paper, we focus on a single electron energy of 1 MeV . Accelerator design efforts and science goals of the beam injection experiment have converged on 1 MeV as the target electron energy. In the sections that follow we discuss how the modeled affects are expected to vary with electron energy, but we do not provide those simulation results in this paper.

We simulate an accelerator design that outputs a pulse of 1 MeV electrons that total 5 J of energy (3.1×10^{13} electrons), and outputs pulses every 5 ms . Neubert and Gilchrist (2002) noted that high currents are required for MeV beams ($>100\text{ A}$) to become unstable; our beam current is only $\sim 1\text{ mA}$. In this work we consider two scenarios: a sequence of 20 pulses

spanning 100 ms and totaling 100 J , and a sequence of 200 pulses spanning 1 s and totaling 1 kJ . In each section that follows, we discuss the effects of increasing or decreasing the total beam energy, or changing the time sequence of pulses. Note that in this paper the “electron energy” refers to the individual electrons (i.e., 1 MeV), while the “pulse energy” or “beam energy” refers to that of the electron pulse (5 J) or sequence of pulses (100 J).

A beam of 1 MeV electrons injected from a distance of $10 R_e$ in a dipole field was simulated by Porazik et al. (2014), who then calculated the spatial, energy, and pitch angle distributions of this beam as it reached 300 km altitude. Those distributions, shown in **Figure 1**, are used as the input distributions to our Monte Carlo modeling. A 2D histogram of the particle positions shows that the beam is distributed approximately as a Gaussian with a 1-sigma radius of 311 m at 300 km altitude. This beam size, together with the pulse energy of 5 J in 5 ms or 1 kJ/s , yields an average flux of about $3 \times 10^5\text{ electrons/cm}^2/\text{s/str}$, comparable to outer radiation belt fluxes at these energies. The beam is extremely field-aligned, with a divergence of <1 degree, due to the careful choice of the firing direction as described in Porazik et al. (2014). However, simulations show that as long as the beam is inside the loss cone, the pitch angle distribution plays only a small role in the atmospheric signatures. For example, a beam with all electrons at 60 degree pitch angle at 300 km altitude, just inside the loss cone, will have a similar energy deposition profile, but raised in altitude by 4 km .

Although ionization, optical, and X-ray signatures scale linearly with the beam energy (for the same electron energy), the choice of electron energy affects these signatures differently. Optical emissions and secondary ionization (leading to electron density disturbances) are nearly proportional to the total energy deposition, as described in section 3. However, X-ray emissions change considerably with electron energy, as the efficiency of bremsstrahlung production increases rapidly at higher energies. As a rule of thumb, approximately 0.2% of the beam energy is converted to X-rays for an electron energy of 1 MeV , while 2% is converted to X-rays for 5 MeV (Krause, 1998).

Ionization production is proportional to the beam energy; we use the rule-of-thumb from Rees (1963) that every 35 eV of energy deposited produces one electron-ion pair. This relationship was validated in Krause (1998). The ionization pair production is then used as a driving source in mesospheric chemistry models, including the Glukhov-Pasko-Inan (GPI) chemistry model (Glukhov et al., 1992; Lehtinen and Inan, 2009) and the Sodankylä Ion and Neutral Chemistry (SIC) model (Verronen et al., 2005; Turunen et al., 2009) to calculate electron density disturbances in the mesosphere and D-region ionosphere, along with the chemical response described in section 4. Here, the response is very strongly dependent on the electron energy. Due to higher electron-neutral collisions at lower altitudes, recombination rates are much higher, and so the electron density perturbation is suppressed. Because higher electron energies deposit energy at lower altitudes, they have a much weaker effect on the electron density disturbance for the same total energy.

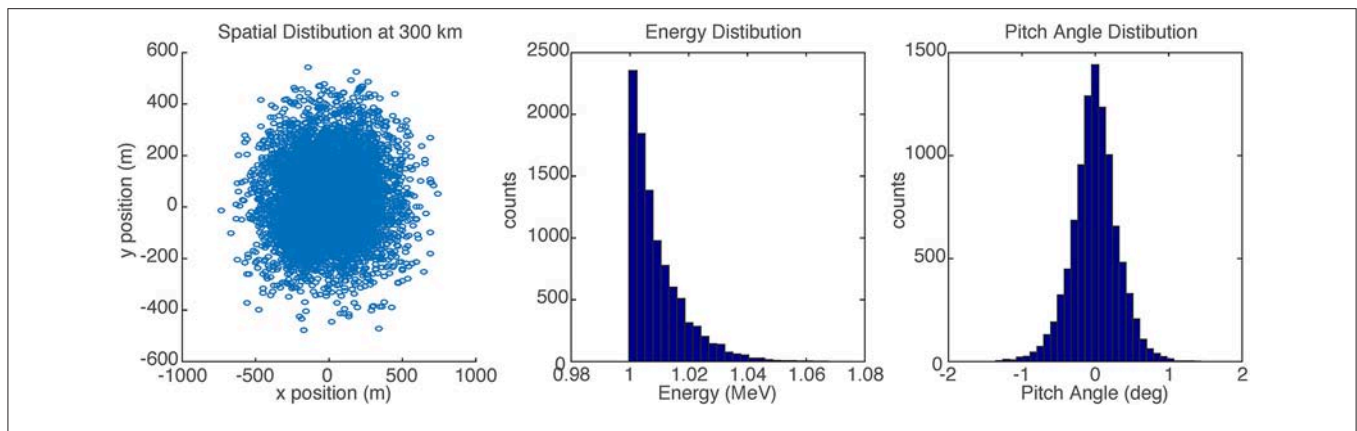


FIGURE 1 | Input spatial, energy, and pitch angle distributions at 300 km altitude for the 1 MeV electron beam, based on simulations from Porazik et al. (2014).

3. IONIZATION AND OPTICAL EMISSIONS

In this section we revisit the ionization and optical signatures that were calculated in Marshall et al. (2014). Here, we use an electron energy of 1 MeV, and beam energies of 100 J or 1 kJ. These beams are actually divided into “pulses” every 5 ms and “subpulses” of 0.5 ms; however, most of the signatures we describe in this paper are not sensitive to the details of the pulse shape on times scales faster than 100 ms. **Figure 2** shows the energy deposition profile for these two beams, along with the optical emission profiles for the 100 J beam. The ionization profile follows that of the energy deposition, under the approximation that every 35 eV deposited creates one electron-ion pair.

We observe from **Figure 2** that the energy deposition scales linearly with the beam energy. The optical emissions are dominated by N_2 first positive (1P) and second positive (2P) emissions, both of which are spread over a large range of wavelengths; as such, detection above the background is much more difficult. The next highest intensity is the N_2^+ first negative (1N) band system, which is heavily concentrated in two bandheads at 3,914 and 4,278 Å. Note that the N_2^+ Meinel (M) band system has a relatively long lifetime ($\sim 6 \mu s$) and a relatively high collisional quenching rate with N_2 , and so is quenched below ~ 90 km altitude where there is appreciable N_2 .

From our calculations of optical emissions, we can generate an energy partitioning that describes the fraction of the total injected energy that is emitted in each of a number of important lines and bands. The results, shown in **Table 1**, show the energy emitted as photons, after accounting for quenching and cascading. This partitioning is consistent with that of typical auroral emissions (Vallance Jones, 1974).

We focus on the N_2 1P emissions, where we can zero-in on narrow emission lines (either 3,914 or 4,278 Å), reducing the background signal. We observe that about 2.2% of the total injected energy is converted to N_2 1P emissions, and 0.6% is converted to N_2^+ 1N emissions. The atomic oxygen green and red line emissions are extremely weak, because the oxygen density is very low at 60 km altitude, and these emissions are rapidly quenched, with 0.7 and 110 s lifetimes.

Using this partitioning table, it is straightforward to make a back-of-the-envelope validation calculation of the expected signal seen by a detector on the ground. We consider an instrument designed to measure the 3,914 Å bandhead of the N_2^+ 1N system. A filter spanning 3,800–3,920 Å will capture 27% of the total 0.6 J emission; assuming a wavelength of 3900 Å, this fraction totals 3.2×10^{17} photons emitted in our band of interest. For simplicity we assume that all photons are emitted from 60 km altitude, and that they are emitted isotropically; and based on MODTRAN (Berk et al., 1987) simulations of the atmospheric transmission, we assume $\sim 40\%$ of the emission reaches the ground, while the rest is scattered or absorbed in the atmosphere. From these values we expect 1.7×10^6 photons/m² over the duration of the beam to reach the ground in our band of interest.

As an approximate instrument response, we consider an optical aperture of 50 mm diameter (a typical camera lens) with a field-of-view that is larger than the emitting region; then we can expect 3.3×10^3 photons to be collected by the instrument. If the instrument is PMT-based, we can consider a window transmission efficiency of 90% and a PMT quantum efficiency (QE) of 28%. We consider instrument dark noise of 2 mA and background airglow of 2 Rayleighs per Ångström (R/Å) at 3,900 Å (Broadfoot and Kendall, 1968). With these noise sources together with shot noise, we calculate an expected signal-to-noise ratio in this PMT instrument of $SNR \simeq 25$ when sampled at 100 Hz.

Instead of a PMT-based system, we also consider measuring the beam spot with a wide field-of-view camera system. In this case, we start from the same 3.3×10^3 photons to be collected by the instrument, assuming the same 50 mm diameter lens. The camera may have the same window transmission of 90%, but a higher QE of 60%, dark current of 0.0003 electrons/pixel/sec, and ~ 1 electron read noise. For such a system averaging frames to 5 fps, we expect an $SNR \simeq 3.6$, assuming the entire beam is contained in a single camera pixel. If instead the beam is spread over a few pixels, the SNR will be reduced from this value.

These calculations show that the optical signature from a 100 J beam should be detectable by either PMT or camera

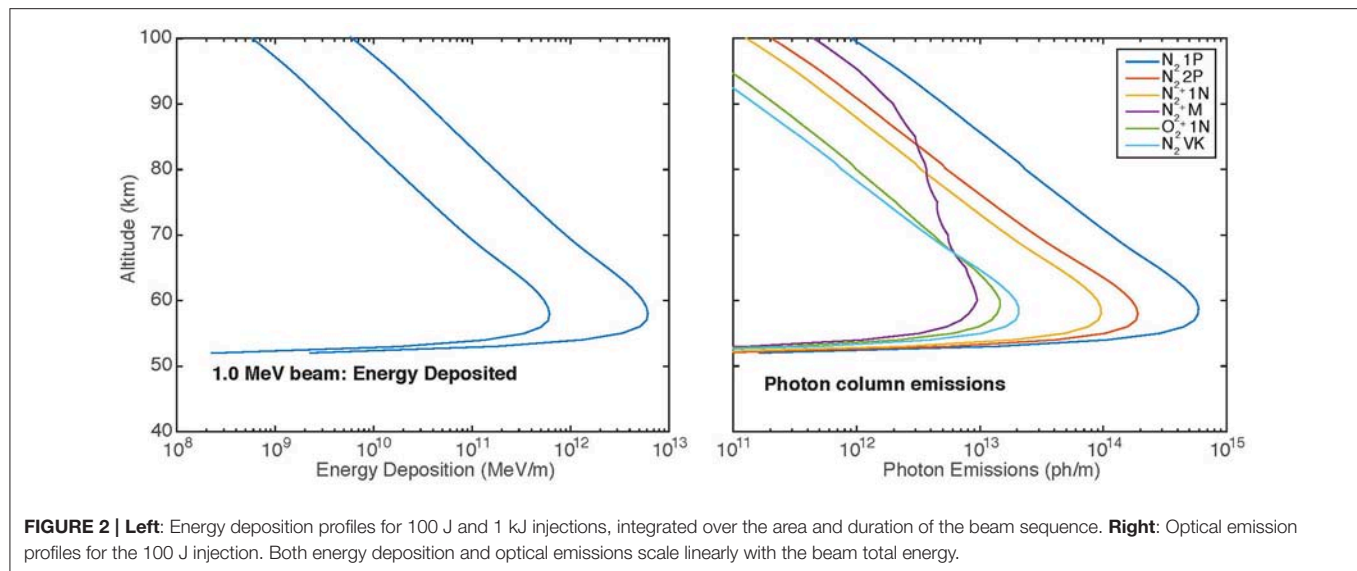


TABLE 1 | Energy emitted in different optical bands from 100 J injected.

Emission	Description	Energy
N ₂ 1P	N ₂ First Positive	2.2 J
N ₂ 2P	N ₂ Second Positive	1.2 J
N ₂ ⁺ 1N	N ₂ ⁺ First Negative	0.6 J
N ₂ ⁺ M	N ₂ ⁺ Meinel	0.1 J
O ₂ ⁺ 1N	O ₂ ⁺ First Negative	0.06 J
O ¹ D	green line, 5577 Å	0.7 μJ
O ¹ S	red line, 6300 Å	0.3 nJ
Total optical energy		<5 J

systems. The PMT system has the advantage of time resolution, valuable if the pulsing sequence is rapid and contains subpulses (but only if the subpulses maintain their separate during the beam propagation). The camera system has the advantage of simple spot detection in a sequence of images, as well as measurement of the spot location, invaluable for field-line tracing applications.

4. CHEMICAL EFFECTS

It is well-known that precipitation of relativistic electrons into the mesosphere can affect the chemistry of this region of the atmosphere. In particular, radiation belt precipitation leads to enhancement of odd nitrogen (NO_x; Rusch et al., 1981) and odd hydrogen (HO_x; Solomon et al., 1982), which ultimately affect ozone concentrations. In this section, we wish to investigate the possible chemical impact of our relativistic electron beam on the upper atmosphere.

The GPI model is a five-species model that includes electrons, heavy and light positive ions, and heavy and light negative ions; as such it cannot calculate the response of individual constituents of interest, such as NO, NO₂, and so forth. Instead, we use the

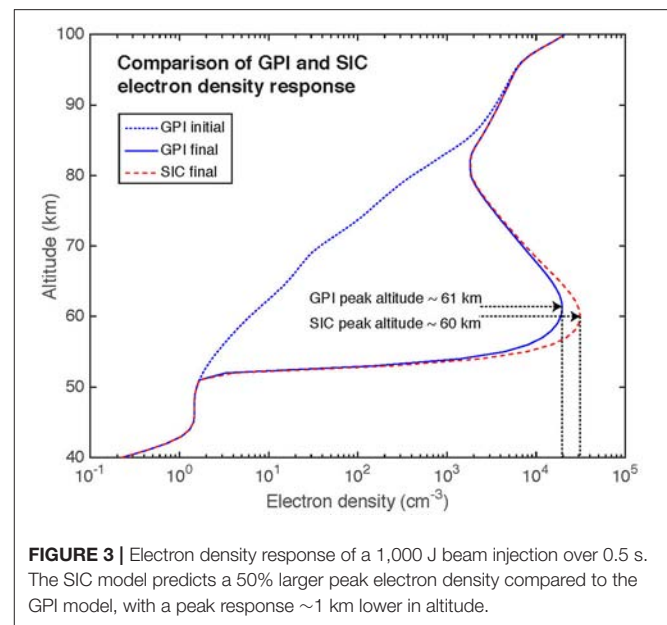


FIGURE 3 | Electron density response of a 1,000 J beam injection over 0.5 s. The SIC model predicts a 50% larger peak electron density compared to the GPI model, with a peak response ~1 km lower in altitude.

SIC model to calculate the response of these species. Described in detail in Verronen et al. (2016), SIC includes forcing from solar UV and soft X-rays, electron and proton precipitation, and galactic cosmic rays. The model solves for the densities of electrons and 70 ions, of which 41 are positive and 29 negative, and 34 neutral species, including O and O₃; N, NO, NO₂, NO₃, and other species cumulatively referred to as NO_x; and H, OH, HO₂, and H₂O₂, cumulatively referred to as HO_x.

In Figure 3, we compare these two models directly, using a 1 kJ beam injection over 100 pulses spaced every 5 ms. The two models use the same initial, background electron density, and then calculate the time response of the electron density profile. The two models compare favorably, though the SIC model

predicts a 55% higher peak electron density, at an altitude 1 km lower than the GPI model. Considering the simplifications made in the GPI model to limit it to five species, this comparison shows that the GPI model provides a reasonably accurate estimate of the electron density response.

Figure 4 shows the time-resolved electron density evolution along the beam axis for the 100 J (20 pulses in 100 ms) or 1 kJ (200 pulses in 1 s) beams, including the electron density profiles after each pulse. In both cases we observe that the electron density begins to saturate at the lower altitudes, due to the rapid recombination at these lower altitudes. After 20 pulses, the peak electron density of $3.9 \times 10^9 \text{ cm}^{-3}$ occurs at an altitude of 59 km. After 200 pulses, the peak electron density is $1.2 \times 10^{10} \text{ cm}^{-3}$ at 63 km altitude.

To determine the chemical response of the atmosphere, and in particular the NO_x , HO_x , and ozone signatures, we use the SIC model. **Figure 5** shows the relative disturbances along the beam axis to each of these species after the 1 kJ injection over 0.5 s. The NO_x density increases by only 0.5% from its background density, with a peak near 70 km altitude. The HO_x density increases by about 0.4%, with the peak at 58 km. The ozone signature is negligible. Therefore, none of the beam applications under consideration should produce any deleterious side effects on the atmosphere.

This small chemical response is encouraging, as it shows that this artificial beam injection will not have a significant, lasting effect on the atmosphere. Energetic electron precipitation is known to produce enhancements in NO_x and HO_x and the former can destroy ozone in the stratosphere. The ionization signature of our electron beam exceeds that of a typical radiation belt electron precipitation event; however, because the spatial extent is very small, but more importantly because the time duration of this precipitation is so short, the effect on atmospheric chemistry is negligible.

The chemical response shown in **Figure 5** is for a single beam pulse, injecting 1 kJ of energy in 0.5 s. While the chemical

response of this pulse is very small, it is possible that a sequence of pulses in the same region of the atmosphere could have a cumulative effect that is more pronounced. Seppälä et al. (2018) modeled the chemical response to a series of microbursts, with comparable time duration and density to our beam pulses, and showed a significant cumulative effect of enhanced NO_x and HO_x . However, their results considered a series of microbursts over a 6-h duration. Microbursts and microburst regions are also likely to cover a much larger spatial scale than our $<1 \text{ km}$ electron beam (Blake and O'Brien, 2016; Crew et al., 2016); as such the beam experiment is unlikely to be able to produce a significant number of pulses in the same region of the atmosphere.

5. ELECTRODYNAMICS AND SPRITE TRIGGERING

Some of the first work on artificial relativistic beam injection was conducted by Banks et al. (1987), Banks et al. (1990), Neubert and Banks (1992), and Neubert et al. (1996). Soon after, in the early 2000s, research into upper atmospheric discharges known as sprites was reaching maturity (e.g., Neubert et al., 2005; Inan et al., 2010). Neubert and Gilchrist (2004) went on to investigate the beam effects in the atmosphere, and suggested the possibility that the relativistic electron beam, upon its interaction with the atmosphere, could modify the conductivity enough to enhance the triggering of sprites at their typical triggering altitude of $\sim 75 \text{ km}$ (Stenbaek-Nielsen et al., 2010; Pasko et al., 2012). Here, we quantitatively assess that possibility using electrostatic field simulations.

We simulate the electron density disturbance in the upper atmosphere as described above and shown in **Figure 4**. This disturbance is three-dimensional in nature, based on the beam spreading calculated in the Monte Carlo model, and is approximately Gaussian with a radius of $\sim 300 \text{ m}$. Next, we

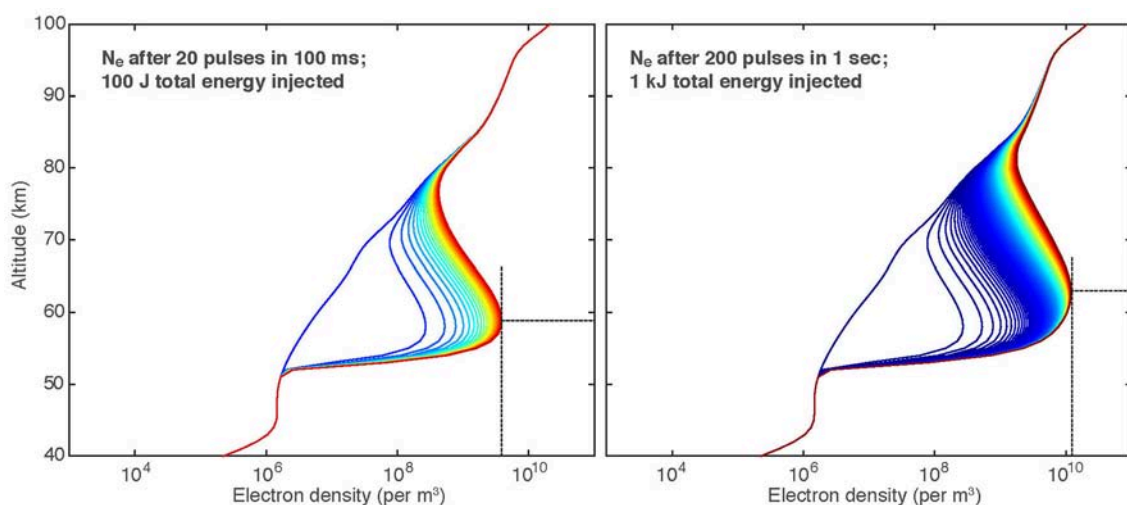


FIGURE 4 | Electron density vs. altitude for a 100 J beam injection (left) and a 1 kJ beam injection (right). Blue to red colors show the electron density after each pulse. Dashed lines mark the peak density and altitude.

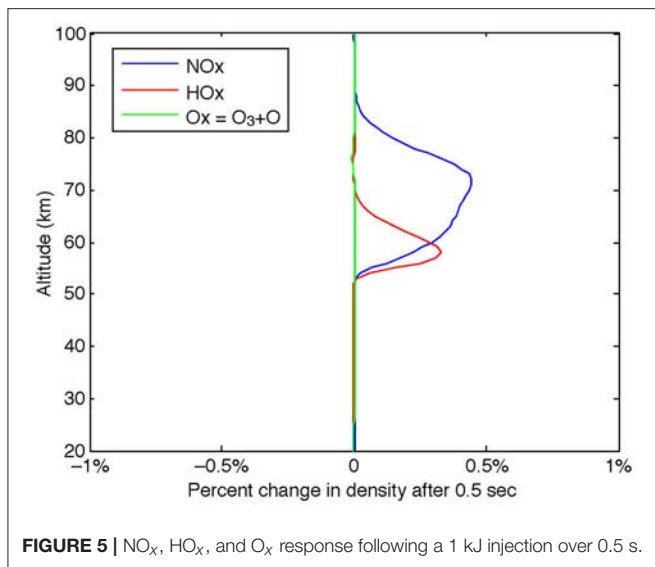


FIGURE 5 | NO_x, HO_x, and O_x response following a 1 kJ injection over 0.5 s.

insert this electron density disturbance into the 2D, cylindrically-symmetric quasi-electrostatic (QES) field model of Kabirzadeh et al. (2015, 2017) and calculate the resulting electric fields. The model is quasi-electrostatic because it solves for dynamically-changing electric fields as time-changing driving sources (charge and current densities) are included.

By default, the QES model uses a uniform grid with either 500 or 1,000 m spatial resolution; however this resolution is clearly insufficient to resolve our 300 m radius disturbance. In order to avoid an excessively large simulation space, the model was modified to use a non-uniform grid; the horizontal resolution is ~ 70 cm at the beam axis, and smoothly increases non-linearly to the maximum grid size of ~ 350 m at a distance of 100 km. The grid is uniform with 250 m resolution in altitude, extending to a maximum altitude of 100 km.

The simulation uses the same background and perturbed ionosphere profile as in **Figure 4**. The electron-neutral collision frequency profile is determined using the method described by Marshall (2014). We wish to determine how the beam injection will change the electric field structure above a thunderstorm. To this end, we calculate the electric fields following the removal of 50 C of charge in a cloud-to-ground lightning discharge. Initially, a -50 C charge is placed at 5 km altitude, and a $+50$ C charge is placed at 10 km altitude. The uppermost charge is removed from the cloud (a 500 C-km charge moment change), causing the electrostatic fields to reconfigure.

The QES model solves for the time-resolved electric fields, but **Figure 6** shows the fields after they have settled; the simulation does not account for charge reconfiguration in the cloud after the discharge. The top row shows the fields without a beam injection, while the bottom row shows the fields after the 1 kJ beam injection. The first two panels in each row show the horizontal (E_r) and vertical (E_z) field components. The rightmost panel shows the reduced electric field E/E_k , i.e., the field magnitude normalized by the breakdown field, which scales with atmospheric pressure. For a sprite to initiate, we are looking for $E/E_k \geq 1$, or $\log_{10}(E/E_k) \geq 0$.

Note that in this scenario, the beam is injected directly above the lightning discharge, where E/E_k is maximum, and that the beam is injected immediately following the lightning discharge, so that the beam modifies the post-discharge field configuration.

Figure 7 shows 1D slices of the reduced electric field (E/E_k) along the beam axis and at ranges from one to 50 km. These are the electric fields immediately after the discharge described above; these fields will recover back to the ambient fields in tens of seconds (e.g., Inan et al., 1996). Following the beam injection we observe significant variation in the field structure, especially on the beam axis at ~ 55 km and between 75 and 85 km. At 55 km the field is perturbed around the bottom of the electron density column. At the higher altitudes, as shown in the zoomed-in panel, we see that $E > E_k$ within 1 km of the beam axis, while $E < E_k$ before the beam injection. These results show that such an experiment could be made to increase the high-altitude electric fields and potentially trigger sprites, but only with very careful timing and fortuitous location.

These results provide an indication that the electron beam may be able to enhance the electrostatic field at sprite altitudes enough to trigger a discharge. However, we have not included the effect of the Earth's magnetic field, which at mid-latitudes, where lightning occurs, is strongly inclined. The magnetic field will push the ionization profile to slightly higher altitudes, and in turn affect this discharge triggering.

6. RADAR AND VLF SCATTERING SIGNATURES

Similar to section 4 above, ionization production in pairs/m³/s are used as an input to chemistry models to determine the expected response of the mesosphere. We use both the SIC model described in section 4 as well as the GPI chemistry model. The latter is a four-species (Glukhov et al., 1992) or five-species (Lehtinen and Inan, 2009) simplified 1D model of mesospheric chemistry that considers electrons, light and heavy positive ions, and light and heavy negative ions. The set of ordinary differential equations relating the densities of these five species are presented in Lehtinen and Inan (2009). The modified electron density profile is used in this section to determine the expected radar and VLF scattering signatures, if any, that could be observed using these techniques.

6.1. Radar Scattering

Using our ionization profiles in section 3, we calculate the time-resolved electron density response in the mesosphere to determine the peak electron density expected as well as the recovery time of this signature. **Figure 8** shows the resulting electron density disturbance and its evolution with time, for the 1 s duration of the pulse sequence (total energy of 1 kJ) and 1 s of its recovery. The background ionosphere density profile is calculated by the SIC model simulations; however the electron density disturbance is so large that the choice of background profile is not important. As in Marshall et al. (2014), the electron

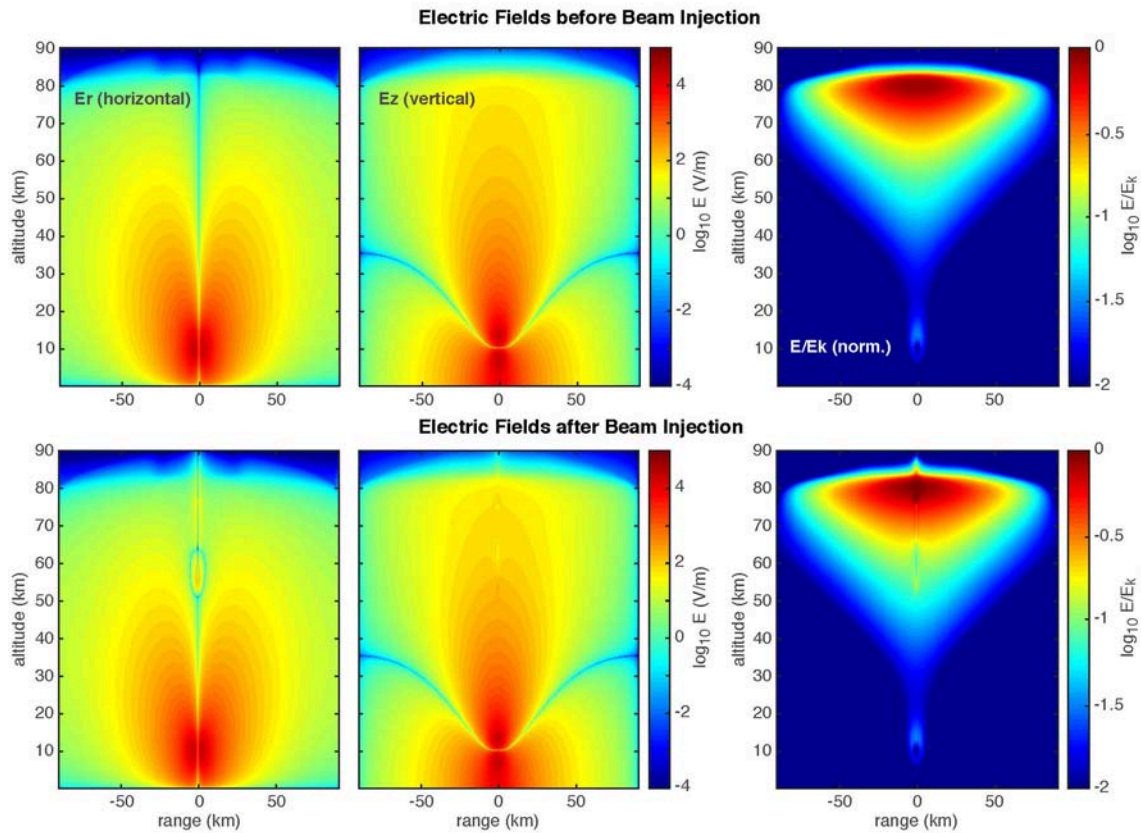


FIGURE 6 | Electrostatic fields (**top**) before and (**bottom**) after beam injection. Rightmost panels show the normalized field E/E_k on a $\log_{10}$ scale.

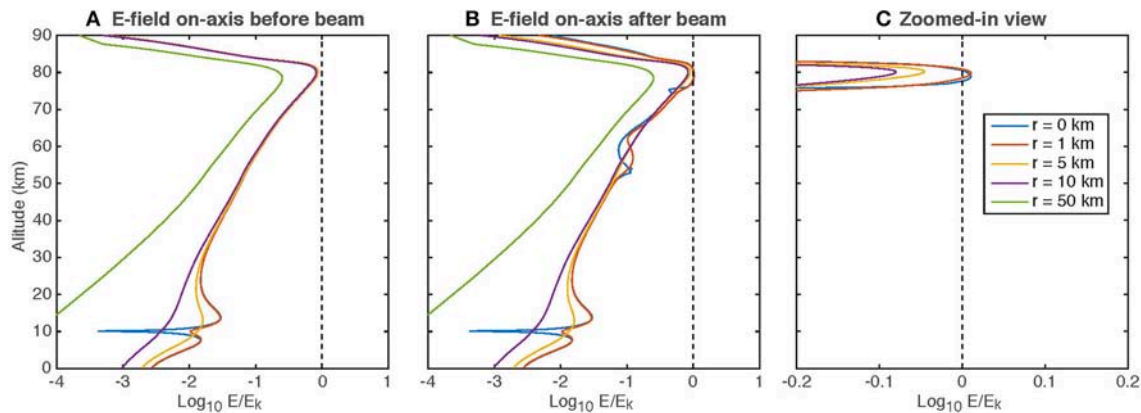


FIGURE 7 | Normalized electric field $\log_{10}(E/E_k)$ in 1D slices along the beam axis and at different radii. **Left:** Before the beam injection; **middle:** after beam injection; **Right:** zoomed-in view after beam injection.

density disturbance recovers back to the background profile over timescales from tens of ms to many seconds, depending on altitude.

We consider the Poker Flat Incoherent Scatter Radar (PFISR) for our detectability estimate. We convert the electron density profile into an expected radar signal-to-noise ratio (SNR) using

the relationship:

$$\text{SNR} = 3.5 \times 10^{-12} \left(\frac{r_0}{r} \right)^2 \frac{2n_e}{(1 + k^2\lambda_D^2)(1 + k^2\lambda_D^2 + T_r)} \quad (1)$$

where n_e is electron density, $r_0 = 100$ km, r is range (or altitude if the radar beam is pointed toward the zenith), $k = 4\pi f/c$ is

the Bragg wavenumber for the radar, $\lambda_D = \sqrt{\epsilon_0 k_B T / q_e^2 n_e}$ is the electron Debye length in the plasma, and $T_r = T_e / T_i$ is the ratio of electron and ion temperatures. Thus the radar SNR is a function of the electron density and both the electron and ion temperatures. For our forward calculations of SNR, we assume $T_r = 1$ in the lower D-region, as electrons and ions are well-thermalized to the neutral temperature through the high collision frequency. At high electron density, $k\lambda_D \ll 1$ and the equation simplifies to $\text{SNR} = 3.5 \times 10^{-12} n_e (r_0/r)^2$; however the full relationship is required below ~ 80 km where $k\lambda_D > 1$. The relationship in Equation (1) was derived for a PFISR experiment using 130 μs , 13-baud Barker codes; the expected SNR will change for different pulse lengths and radar performance.

A radar SNR < 1 does not mean the signal cannot be detected; by averaging pulses coherently and incoherently, we can dramatically improve the detectability, which is estimated as follows. We can combine N consecutive radar pulses using coherent averaging, up to the correlation time, which is about 200 ms at 65 km altitude. For a Lorentzian radar spectrum, the coherent processing gain is given by

$$G = \frac{1}{N} \sum_{n=0}^{N-1} \sum_{m=0}^{N-1} e^{-2\pi\omega|m-n|t} \quad (2)$$

where ω is the spectral width in Hz (about 5 Hz at 65 km altitude), and t is the inter-pulse period, taken to be 2 ms for this experiment. (Note that the radar “pulses” here, every 2 ms, are not the same as the beam “pulses” transmitted every 5 ms.) With $N = 100$ radar pulses in 200 ms, we calculated a coherent processing gain of $G = 27$. Finally, the relative error in the ISR power estimate is found by incoherently averaging K sets of coherent averages (e.g., Farley, 1969):

$$\frac{dS}{S} = \frac{1}{\sqrt{K}} \left(1 + \frac{1}{G \cdot \text{SNR}} \right) \quad (3)$$

where K is the number of incoherent averages, in this case $K = 5$ to represent the five 200 ms periods. This relative error is plotted in the right panel of **Figure 8**. A relative error of $dS/S = 1$ indicates that the signal is 1σ above zero SNR; $dS/S = 0.5$ indicates 2σ above zero SNR; $dS/S = 0.33$ indicates 3σ above zero SNR, and so forth. We observe that the maximum SNR in these results is about -10 dB, corresponding to a minimum $dS/S = 0.27$. This shows that the electron beam pulse sequence of 1 kJ injected over 1 s should be marginally detectable by an incoherent scatter radar such as PFISR, when it is running Barker-type emission codes. Newer codes that increase the coherent gain combined with longer integration times will decrease dS/S and thus improve detectability. For example, a 50% increase in the averaging intervals would decrease dS/S to 0.22. Similarly, the electron beam signatures will likely be observable by the upcoming EISCAT 3D radar (Turunen, 2009).

Note that the radar signal is not sensitive to the background state of the ionosphere below 80 km; PFISR sees only noise below these altitudes under typical conditions, irrespective of the background D-region ionosphere state. The exception only occurs under very intense radiation belt precipitation, which can

be detectable by PFISR below 70 km, and which may interfere with our beam detection.

6.2. Subionospheric VLF Scattering

Next, we consider whether or not the electron density disturbance from the beam would be observable through scattering of subionospheric very-low-frequency (VLF) transmitter signals. Powerful ground-based transmitters operated by the US Navy radiate VLF waves into the Earth-ionosphere waveguide, and the amplitude and phase of the signals observed by a distant receiver are particularly sensitive to the D-region ionosphere. At night, these waves reflect from altitudes ~ 80 – 85 km and are modified by electron density disturbances below the reflection height.

We test the possibility that the beam will perturb the VLF signal by simulating the propagation of VLF transmitter signals in the Earth-ionosphere waveguide, and comparing the amplitude and phase at a distant receiver before and after the beam pulse. We use the Finite-Difference Time Domain (FDTD) propagation model of Marshall (2012) and Marshall et al. (2017), which allows calculation of amplitude and phase for any frequency at any distance, and allows for small-scale ionospheric disturbances with ~ 500 m resolution or better.

The simulations shown here use a grid resolution of 500 m; as such, the electron density disturbance created by the beam only spans a few grid cells. We simulate a scenario shown in **Figure 9**, using the NLK transmitter. The pulse is injected above Poker Flat, AK, and a receiver is located 500 km further along the great-circle-path (GCP) connecting NLK and Poker Flat. To reduce the simulation time, we use a virtual transmitter 1,000 km away from Poker Flat instead of simulating the entire path. **Figure 9** also shows the final electron density along the simulation path, with the beam injection shown at 1,000 km from the transmitter.

To estimate the expected amplitude and phase perturbations to the VLF signal, we run two simulations: one without the beam-induced disturbance, and one with the disturbance. **Figure 10** shows the amplitude and phase along the ground for both cases; the rightmost panels are zoomed-in views of the last 500 km. We see that the VLF signal is significantly perturbed, with up to 1 dB of amplitude change and ~ 10 degrees of phase change. For reference, in comparable VLF data, a minimum detectable perturbation is about 0.1 dB and 1 degree. Note that the ringing at the end of both simulations is a numerical artifact, due to the simulation being stopped before the highest-order modes have equilibrated.

The natural variation in the D-region ionosphere will lead to variation in the received VLF signal amplitude and phase, as well as the received perturbation. The D-region variations can lead to amplitude variations at night of $\sim \pm 5$ dB, and phase variations of $\sim \pm 50$ degrees, varying on time scales of minutes to hours. A more comprehensive study, left to future work, is needed to assess the expected VLF perturbation under changing D-region conditions.

It is tempting to conclude that these results show that the subionospheric VLF method may be able to detect the beam-induced electron density disturbance, but we cannot yet make this conclusion. These simulations are 2D only, in range and

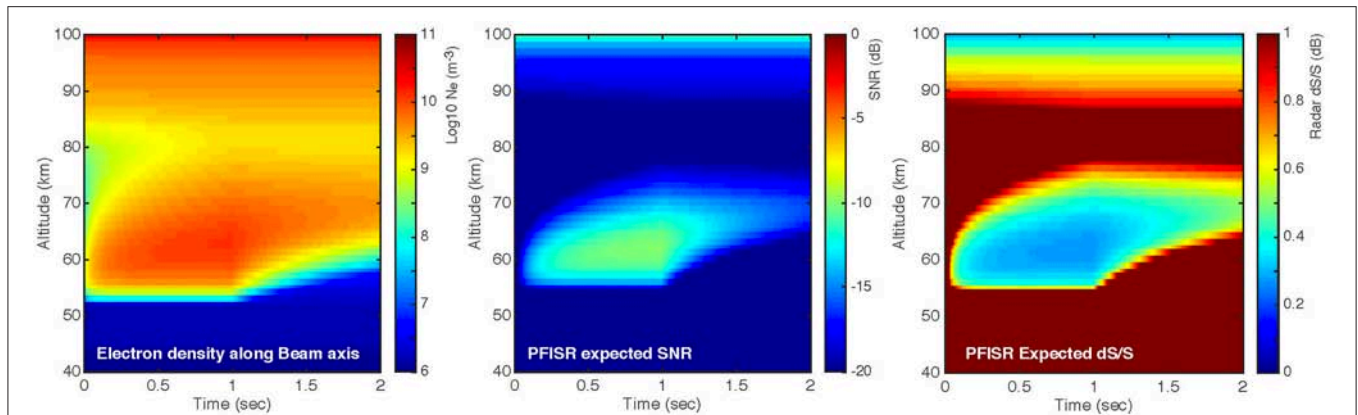


FIGURE 8 | Left: Time-resolved electron density response for 1 s beam injection and 1 s recovery. **Middle:** expected SNR using PFISR radar parameters. **Right:** expected relative error, dS/S , again using PFISR radar parameters and pulse averaging as described in the text.

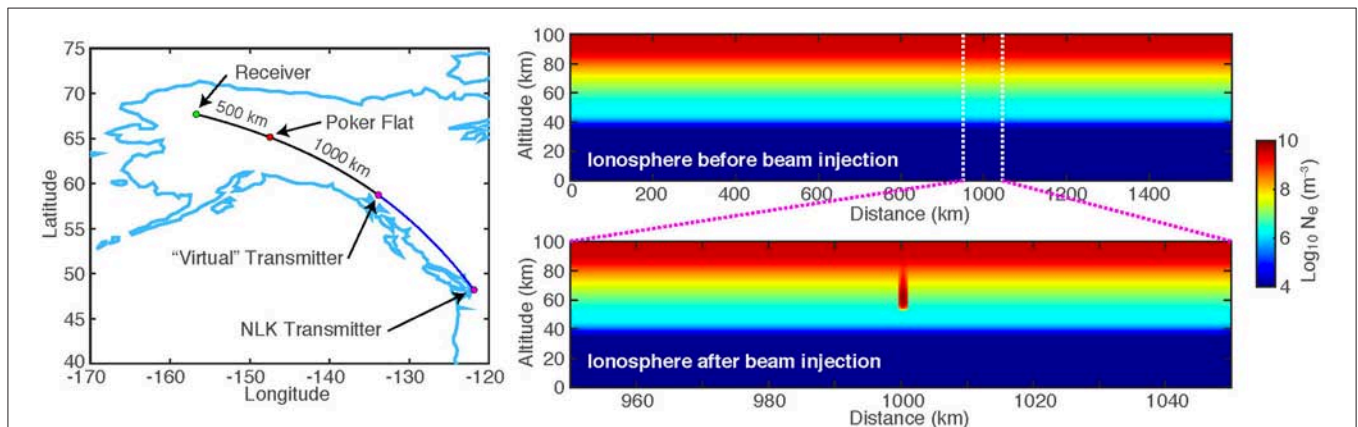


FIGURE 9 | VLF scattering setup. To simulate the NLK transmitter we use a virtual transmitter 1,000 km from Poker Flat along the same great-circle-path. The receiver is 500 km past Poker Flat. **Right:** electron density profile along the path, showing the beam disturbance at 1000 km range from the transmitter.

altitude. As such, the disturbance imposed is effectively infinite in the third dimension; rather than a column of electron density, we have imposed a “wall” extending in and out of the page. This configuration is likely to produce a larger scattered signature than a single 300 m radius column. To better quantify the expected amplitude and phase perturbation, a full 3D simulation is needed. Such a simulation is extremely computationally expensive, and a single model to make this estimate does not currently exist. Nonetheless, these preliminary 2D simulations do not rule out the possibility that the subionospheric VLF method may be able to detect the beam-induced disturbance.

7. DISCUSSION

A relativistic beam of electrons injected from high altitudes has great potential for field line mapping, wave-particle interaction studies, and atmospheric studies (e.g., Delzanno et al., 2016), but most studies will require detection of the beam in the atmosphere. In this paper, we have provided results of simulations of the interaction of a beam of 1 MeV electrons

with the upper atmosphere in order to assess its detectability via numerous diagnostic techniques. We have further explored the effects of the electron beam on the atmosphere, in terms of the chemical and electrodynamic response of the region. For the latter, we show that the beam injection into the atmosphere may aid in the triggering of sprites at high-altitude, though the inclination of the Earth’s magnetic field must be taken into account.

We simulate a beam of 1 MeV electrons totaling 100 J or 1 kJ of energy. Monte Carlo simulations provide an estimate of the ionization produced by these beams and the altitude distribution and horizontal distribution of this ionization. Optical emissions are then calculated from the ionization production, and we determine the photon production taking into account quenching and cascading in a suite of N_2 , N_2^+ , and O_2^+ emission band systems. We estimate the expected signal-to-noise ratios in a photomultiplier tube (PMT)-based detection system, and in an all-sky camera. These two systems have different advantages. The PMT system has the speed (1 kHz) necessary to detect individual sub-pulses in the beam pattern, but does not have any spatial

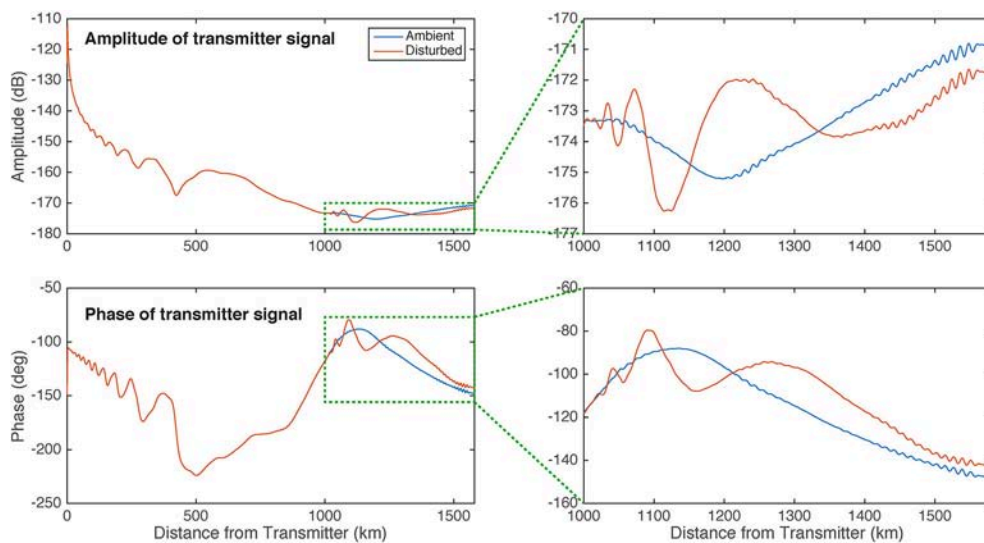


FIGURE 10 | VLF amplitude and phase along the ground with and without the beam injection. A VLF receiver at 1,500 km would be expected to measure an amplitude change of ~ 1 dB and a phase change of ~ 10 degrees.

information; the all-sky camera system lacks time resolution but can locate the beam spot in the sky with high accuracy, a critical requirement for many of the science applications of the electron beam.

Both systems should have sufficiently high SNR to detect the spot in the upper atmosphere. However the SNR values calculated in section 3 depend on a number of parameters which will vary for different systems. In particular, the PMT system depends on the choice of PMT and its wavelength response, noise characteristics, and the instrument sampling rate, in addition to optical design parameters. The camera system similarly depends on parameters of the camera chosen and the optical system, including filter transmission and passband. As such, the expected SNR will vary for different systems, and the system must be carefully designed to be optimized for the expected signatures. However, our calculations of the SNR and detectability are validated by the SEPAC experiments (Neubert et al., 1995) who observed optical emissions of 1–5 kR with a factor of 7.5 higher energy flux than our proposed experiments.

Similarly, the radar signatures presented in section 6.1 must be considered for a particular experiment design. The SNR calculated by Equation (1) will change for different radar pulse parameters and for different ISRs. What's more, the detectability in ISR appears to be marginal with standard radar beam pulse codes. New beam schemes and longer integration times will be required to ensure detectability.

One important characteristic of ISRs that must not be overlooked is that these instruments can provide detection of the electron beam in dayside conditions, where optical detection methods are not possible. An entire class of magnetospheric phenomena, such as plasma entry through the dayside boundary layer between the solar wind and the magnetosphere, still have several fundamental outstanding questions that can be answered with the appropriate match between magnetospheric

in-situ measurements and the unambiguous identification of the ionospheric foot-point.

Subionospheric remote sensing has its own set of difficulties for detection. The receiver needs to be downstream of the ionization patch relative to the transmitter, although some deviation is likely acceptable; the forward scattering of the ~ 300 m radius patch will have some angular distribution. However, it is unclear at this point how strong the scattering will be in a full, three-dimensional scattering scenario. We require a full 3D scattering model to fully assess the VLF scattering. However, such a 3D model will be computationally expensive, since it requires ~ 100 m resolution around the ionization patch, but a transmission distance of thousands of kilometers.

In summary, in this paper we have presented a range of signatures of the 1 MeV beam interaction with the upper atmosphere, and quantified the expected signatures in different diagnostics. Optical detection of the beam spot remains the most promising method, and a combination of PMT and camera detection would allow both time resolution and spatial location of the spot. Radar and VLF detection of the ionization patch are likely marginal, though the latter requires further study.

DATA AVAILABILITY

The simulation outputs used in the analysis and results herein are available for download from Github at <https://github.com/ram80unit/FrontiersBeamPaper>.

AUTHOR CONTRIBUTIONS

RM conducted GPI chemistry simulations and VLF propagation simulations, analyzed simulation outputs, created figures, and

wrote the manuscript text. WX conducted Monte Carlo simulations of electron precipitation in the atmosphere. AK conducted SIC model simulations of the atmospheric response. RK developed the 3D QES model and conducted simulations using that model. ES is the project PI and contributed to the analysis of radar signatures.

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Enhanced Energetic Neutral Atom Imaging

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Over the past two decades, instruments designed to image plasmas in energetic neutral atom (ENA) emission have flown in space. In contrast to typical satellite-based *in situ* instruments, ENA imagers provide a global view of the magnetosphere because they remotely measure ion distributions via neutrals that are not tied to the magnetic field. An intrinsic challenge that arises during analysis of magnetospheric ENA images is that the ENA fluxes are integrated along the line-of-sight of the instrument. We propose a method of enhancing ENA emission from a localized region in space, thereby enabling spatially resolved measurements of ENA emission in a remotely obtained ENA image. Here we show that releases of modest volumes ($\sim 1.4 \text{ m}^3$) of liquid hydrogen in space are sufficient to accomplish the ENA localization.

Keywords: magnetosphere, active space experiment, energetic neutral atoms, charge exchange, geocorona, ion energy spectra

INTRODUCTION

In the late 1980's, Roelof realized that energetic particle signals obtained during cusp transits by the IMP 7/8 and ISEE-1 spacecraft could be explained if the signals were actually energetic neutral atoms that had escaped the inner magnetosphere after charge exchange collisions with energetic ions (Roelof et al., 1985; Roelof, 1987). After he demonstrated that the measurements could be used to create neutral atom images of the inner magnetosphere, the space plasma community embraced the concept of energetic neutral atom imaging. A remote sensing technique, energetic neutral atom (ENA) imaging provides global views of magnetospheric ion populations that have been processed through charge exchange collisions with the Earth's cold geocorona (Scime and Zaniewski, 2004). Where the geocorona coexists with hot plasma, fast ions undergo charge exchange collisions with geocoronal neutral atoms; producing a cold ion and an energetic neutral atom with the same energy as the original ion. By its very nature, the ENA production process yields a flux of ENAs along every line-of-sight through the magnetosphere. Since the cross section for charge exchange collisions between ions and neutrals is well known from laboratory measurements, with a model of the shape and density of the geocorona it is possible to convert energy-resolved ENA measurements into a measure of the line-integrated ion energy spectrum along a given line-of-sight.

For readers interested in an in-depth discussion of neutral atom imaging, Gruntman's review of the history of neutral atom imaging is an excellent resource (Gruntman, 1997). The first ENA imager of the "modern" era, the Ion Neutral Camera (INCA), successfully obtained ENA images of the Saturnian magnetosphere from aboard the Cassini spacecraft (Mitchell et al., 2000). The instrument complement of the Imager for Magnetopause-to-Aurora Global Exploration (IMAGE)

spacecraft included three different ENA cameras. The Low [LENA; (Moore et al., 2000)], Medium [MENA; (Pollock et al., 2000)], and High [HENA; (Mitchell et al., 2000)] energy neutral atom imagers. The energy range of those instruments spanned 15 eV to 500 keV per nucleon. Using the same medium energy instrument design as IMAGE, the two ENA imagers of the TWINS mission (McComas et al., 2009a) have provided nearly a decade of continuous ENA observations of the terrestrial magnetosphere (Keesee and Scime, 2015). Using a different technique to image medium energy neutrals created at the heliospheric termination shock, the IBEX-HI ENA imager aboard the Interstellar Boundary Explorer (IBEX) spacecraft has revolutionized our understanding of the structure of the outer heliosphere (McComas et al., 2009b).

ENA imaging of the magnetosphere has also provided a wealth of information about ion dynamics during geomagnetic storms and substorms. Roelof found a strong day-night asymmetry in the ion ring current during a geomagnetic storm using ISEE 1 ENA data (Roelof, 1987). Pollock used MENA data to demonstrate the evolution from partial to complete ring current during a storm as well as that the loss of ring current ions is dominantly through the dayside magnetopause (Pollock et al., 2001). C:son Brandt used HENA data to show the existence of strong, skewed equatorial electric fields in the inner magnetosphere that depend upon the solar wind velocity and interplanetary magnetic field (IMF) B_y (C:son Brandt et al., 2002b). Keesee et al. (2014) used TWINS images to show the propagation of regions of energized ions from the magnetotail toward the inner magnetosphere, with downward deflection near geosynchronous orbit. Perez et al. (2016) discovered the existence of two ion flux peaks in the inner magnetosphere trapped population during a storm using TWINS data. ENA data have also been used in conjunction with modeling to study the inner magnetosphere electric field (Buzulukova et al., 2010; Fok et al., 2010), the neutral geocorona (Ilie et al., 2013a), ion composition (Ilie et al., 2013b), and the influence of boundary conditions on simulation results (Elfritz et al., 2014).

While ENA imaging has yielded important insights into the structure and evolution of the terrestrial magnetosphere, the lack of spatial resolution has held back widespread acceptance of the ion energy distribution results from ENA energy spectra measurements—even though comparative studies between ENA energy spectra and local ion energy spectra have demonstrated impressive consistency between the two techniques (Keesee et al., 2012, 2014; Perez et al., 2016; Goldstein et al., 2017). Global ENA measurements have also proved useful for model validation (e.g., Fok et al., 2014). Spatially resolved ENA measurements would provide an additional method of using ENA measurements to validate magnetospheric models while increasing the intrinsic value of ENA data.

Therefore, in the spirit of the 2017 workshop on Active Experiments in Space held in Santa Fe, NM, here we describe an active space experiment that would provide spatially resolved ENA measurements of ion energy distributions and enable researchers to distinguish signals from a specific spatial location from within the line-integrated measurement.

ACTIVE ENA MEASUREMENTS

An ENA is created through a charge exchange collision between an energetic ion and a cold neutral atom. The measured ENA intensity, j_{ENA} (with units of $(\text{cm}^2 \text{ sr s eV})^{-1}$), depends upon the ion intensity, j_{ion} , the charge exchange cross section, σ_{cx} (Freeman and Jones, 1974), and the neutral density, n_n , through the relation

$$j_{\text{ENA}}(\mathbf{E}, \vec{\mathbf{u}}) = \sigma_{\text{cx}}(\mathbf{E}) \int_0^{\vec{\mathbf{R}}} \left(n_n(\vec{\mathbf{r}}(s)) j_{\text{ion}}(\vec{\mathbf{r}}(s), \mathbf{E}, \vec{\mathbf{u}}) \exp\left(-\int_0^{\vec{\mathbf{r}}(s)} \left(\alpha(s') ds'\right) d\mathbf{s}\right) \right) \quad (1)$$

where the integral is performed along the line of sight (LOS). Attenuation of ENAs due to additional collisions and photoionization is accounted for in the integral over $\alpha(s')$. We assume this integral is approximately zero, which is applicable to most of the magnetosphere beyond a few Earth radii.

If we assume a Maxwellian parent ion distribution, then the hottest region along the LOS dominates the high energy portion of the spectrum (Hutchinson, 1987). Under this assumption, equation (1) becomes

$$\frac{j_{\text{ENA}}(\mathbf{E}, \vec{\mathbf{u}})}{\sigma_{\text{cx}}(\mathbf{E})\mathbf{E}} \approx \frac{\xi n_n(\mathbf{r}^*) n_i(\mathbf{r}^*)}{\sqrt{2m_i} (\pi T_i(\mathbf{r}^*))^{3/2}} \exp\left(\frac{-\mathbf{E}}{T_i(\mathbf{r}^*)}\right) \quad (2)$$

where the integral has been approximated by the value at the hottest point, $\mathbf{r}^*$, times a characteristic length along r , ξ (Scime and Hokin, 1992). The Maxwellian assumption is valid for the plasma sheet but less so for the inner magnetosphere. The Maxwellian assumption method calculates an average energy of the bulk population that provides information about the energization of the ions within the measured energy range (typically 1–100 keV). The ion temperatures, T_i , are calculated by fitting Equation (2) to the measured ENA energy spectrum, a method that has been verified through comparison to *in situ* measurements (Scime et al., 2002).

The key term in the measured ENA flux for the active method proposed here is the density of the background neutral population, $n_n(\vec{\mathbf{r}}(s))$. The geocorona surrounding the Earth is a roughly spherical distribution of cold neutral gas. Early measurements by Rairden et al. (1986) estimated the neutral hydrogen density at geosynchronous orbit ($6R_E$) to be 50 cm^{-3} . Subsequent measurements by Ostgaard et al. (2003) using extreme ultraviolet (EUV) emission measurements from the IMAGE spacecraft yielded nearly identical values for the geosynchronous neutral density. More recent measurements using EUV measurements from the TWINS mission (Bailey and Gruntman, 2011) yield daytime geosynchronous neutral densities of 100 cm^{-3} . At night, the neutral density doubles to 200 cm^{-3} (Bailey and Gruntman, 2011). ENA images of the magnetotail suggest that the neutral density in the anti-sunward tail of the geocorona decreases more gradually with increasing

distance from the Earth, consistent with a substantial night/day asymmetry in the geocoronal neutral density at geosynchronous orbit (Keesee et al., 2012).

Before describing how the ENA signal might be enhanced for a particular region in space, it is illustrative to describe how ENA images are constructed from the ENA measurements. Three methods are generally used. First, ion distributions can be calculated from the ENA data using a deconvolution technique described by Perez et al. (2000). Second, a constrained linear inversion technique as described by C:son (C:son Brandt et al., 2002a), can be used to find ion distributions that result in ENA images that are matched to the measurements. These two techniques provide a measurement of the ion density and temperature, but require accurate models of the magnetic field and neutral geocorona. The third method, the method typically used in our studies because it requires few computational resources and a minimum of assumptions, enables a calculation of only the ion temperature independent of the neutral source distribution. Typically, we map the measured ENA fluxes along the LOS to the xy plane (GSM coordinates), assuming that the hottest point along the LOS occurs in the central plasma sheet (Hughes, 1995) and, therefore, near the equatorial plane in the magnetotail; which also eliminates the need for a magnetic field model. The field of view (FOV) of each pixel in an ENA image is projected along the LOS to calculate the intersection of the FOV with the GSM xy plane. The ENA emissivity is placed proportionally in the $0.5 \times 0.5 R_E$ xy plane bins with which the FOV intersects. This algorithm accounts for the increasing FOV with increasing distance from the Earth. A modeled magnetosphere boundary (Shue et al., 1997) is used to discard flux projected to bins that are outside the boundary. The average emissivity for each bin is calculated prior to calculating the ion temperature. This method of calculating projected ion temperatures has been validated with *in situ* measurements (Keesee et al., 2008) and a typical ENA ion temperature image is shown in **Figure 1**.

Given that geocoronal hydrogen densities range from 100 to 200 cm^{-3} at $6R_E$ and decrease for larger geocentric distances, an increase in the local neutral hydrogen density of 200 cm^{-3} would produce a significant enhancement of the ENA flux emitted from that region of space. An earlier study when the field of ENA imaging was just beginning (McComas et al., 1993), proposed that cold releases from conventional explosions in space could be detectable with the ENA instruments under development at the time. That study, with important national security implications, is a precursor to the analysis described here.

In terms of a line integrated measurement, the typical lines-of-sight for the TWINS ENA imager (see **Figure 2**) pass through roughly $10R_E$ of the magnetosphere where the plasma and neutral densities are large enough to generate significant ENA fluxes. Assuming uniform ENA emission along the entire line of sight and a uniform geocoronal hydrogen density of 100 cm^{-3} , conservative estimates, an increase of 200 cm^{-3} along $0.5R_E$ would yield a 10% increase in the total ENA signal from a given line of sight. A modulated increase would be easily detectable if

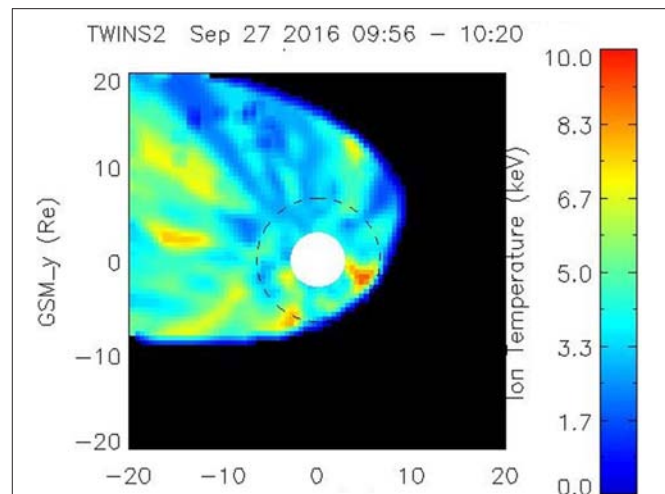


FIGURE 1 | Ion temperatures in the GSM equatorial plane calculated using energetic neutral atom data from TWINS 2 on Sept. 27, 2016 at 9:56–10:20 UT. The white disc has radius $3 R_E$ centered at Earth and the dashed line indicates geosynchronous orbit. The area of measured temperatures is influenced by a combination of the instrument field of view and a modeled magnetosphere boundary.

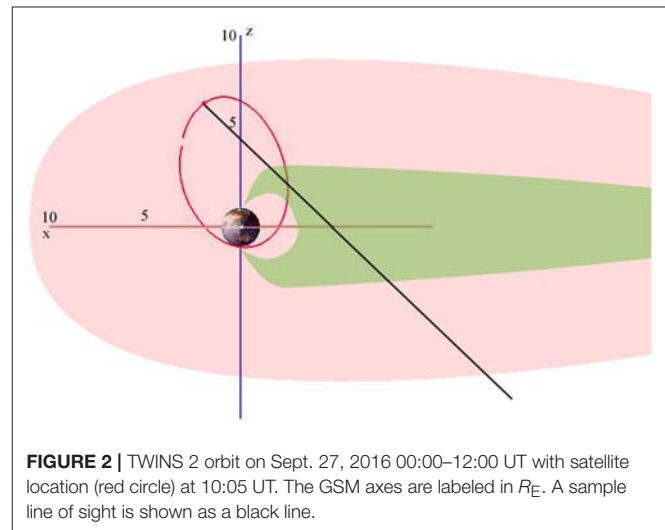


FIGURE 2 | TWINS 2 orbit on Sept. 27, 2016 00:00–12:00 UT with satellite location (red circle) at 10:05 UT. The GSM axes are labeled in R_E . A sample line of sight is shown as a black line.

the modulation frequency was distinct from naturally occurring geophysical processes.

Based on the density of liquid hydrogen ($2.02 \text{ g/mol} = 70.85 \text{ g/l}$), an increase of 200 cm^{-3} in the neutral density across a spherical volume $0.5R_E$ in diameter requires

$$\begin{aligned} \text{atoms of hydrogen} &= 200 \text{ cm}^{-3} \times \frac{4\pi (6.4 \times 10^8 \text{ cm}/2)^3}{3} \\ \text{atoms of hydrogen} &= 2.8 \times 10^{28}. \\ \text{moles of hydrogen} &= 46.5 \times 10^3. \end{aligned} \quad (3)$$

mass of liquid hydrogen (H_2) = 50 kg.

liters of liquid hydrogen (H_2) = 663 liters (0.7 m^3).

This calculation takes advantage of the molecular nature of liquid hydrogen and therefore, to generate neutral hydrogen atoms, a dissociation process would be required, e.g., a UV light source. However, at the energies of interest, ~ 10 keV, the charge exchange cross section for producing ENAs from collisions with molecular hydrogen, 10^{-15} cm², is nearly identical to the cross section for neutral hydrogen (Freeman and Jones, 1974). Therefore, doubling the volume of liquid hydrogen to 1.4 m^{-3} would yield the required increase in the ENA emission. In other words, a cube of liquid hydrogen 1.1 m on a side provides enough molecular hydrogen for the desired increase in the neutral density over a spherical volume $0.5R_E$ in diameter.

To put this volume in perspective, the published fuel capacity of the 2nd stage of a Falcon 9 rocket is 27,634 liters of liquid oxygen and 17,413 liters of liquid kerosene (www.spaceflight101.net/falcon-9-launch-vehicle-information). One such release is therefore equivalent to 7.6% of the standard oxygen capacity by volume of the Falcon 9 2nd stage. Surreptitiously, the charge exchange cross section for protons on molecular oxygen at 10 keV is four times that of protons on molecular hydrogen (Allison, 1958). Therefore, an unaltered Falcon 9 2nd stage boosted into geosynchronous orbit could provide more than 40 controlled releases of liquid oxygen for an active ENA experiment of the type proposed here. There are other options for producing these ENA enhancements as well. For example, unused fuel reserves on standard commercial or scientific launches could be used to trigger the ENA emission enhancements as described here.

Neutral cloud releases are not a novel concept. Fuselier et al. (1994) optically tracked an explosive release of barium to map local magnetic field lines and to investigate the ionization processes in near-Earth space. More recently, barium releases have been proposed as a means of exciting electromagnetic ion cyclotron waves in the lower magnetosphere. Fundamentally, this proposed active experiment is much less complicated than the barium release experiments (Fuselier et al., 1994). The desired material is under pressure and release is accomplished with a standard cryogenic valve.

As the released plume expands, the spatial localization of the ENA enhancement begins. Initially, the energy spectrum of any increases in observed ENA signal can be compared to local measurements of the ion energy spectrum. As the plume expands further, the enhanced ENA emission integrates over a large volume of space and comparisons could be made to ion energy spectra in nearby regions of space. Assuming the released gas is at the boiling temperature of liquid hydrogen, ~ 33 K, the gas cloud will expand at ~ 400 m/s. Comparing that escape speed to the orbital velocity of a geosynchronous spacecraft, 3,100 m/s, it is clear that for most orbital tracks, the motion of a released

cloud will be dominated by the spacecraft velocity and in that frame, a cloud would take roughly 3 h to expand to a radius of $1R_E$. An oxygen cloud would expand much more slowly even if released at the higher temperature of liquid oxygen. Because the expanding cloud will track the spacecraft velocity, the cloud will also sweep through magnetic local time in the magnetosphere as it expands—enabling the sampling of a range of magnetic local times with each release. Another factor in the lifetime of the cloud is the combined photoionization and charge exchange rate at 1 AU. For hydrogen, these two loss processes would reduce the overall cloud density by $<1\%$ over the time it takes for the cloud to expand to $1R_E$ (Ogawa et al., 1995). The lifetime for an expanding oxygen cloud is similar, given the photoionization rate of oxygen at 1 AU (Meier et al., 2007).

While the enhanced ENA emission should be localizable by a single ENA imager, the advantages of stereoscopic ENA imaging, such as that provided by the dual TWINS instruments (McComas et al., 2009a), are significant. Releases such as those proposed here would provide a unique opportunity to validate stereoscopic inversion techniques as well as identify projection issues arising from non-uniform pitch angle distributions of the parent energetic ion populations. An ideal active neutral release experiment might include multiple tanks of liquid hydrogen or oxygen placed into different regions of the magnetosphere. Individual and/or simultaneous enhancements could then be imaged with a single or multiple, full sky, ENA imagers to provide simultaneous measurements of the global ion energy spectrum from different regions of space.

SUMMARY

The next step in validating this active space ENA experimental concept would be to perform computational simulations, using existing magnetospheric models of ENA emission (Perez et al., 2001), to estimate detection thresholds and density enhancement requirements for a variety of magnetospheric conditions and viewing geometries.

AUTHOR CONTRIBUTIONS

ES conceived of the presented idea based on discussions with AK. ES performed the calculations and took the lead on the manuscript. AK contributed to the text and created the figures.

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Ionospheric Non-linear Effects Observed During Very-Long-Distance HF Propagation

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A new super-long-range wave propagation technique was implemented at different High Frequency (HF) heating facilities. The HF waves radiated by a powerful heater were scattered into the ionospheric waveguide by the stimulated field aligned striations. This waveguide was formed in a valley region between the E- and F- layers of the ionosphere. The wave trapping and channeling provide super-long-range propagation of HF heater signals detected at the Ukrainian Antarctic Akademik Vernadsky Station (UAS) which is many thousand kilometers away from the corresponding HF heating facility. This paper aims to study the excitation of the ionospheric waveguide due to the scattering of the HF heating wave by artificial field aligned irregularities. In addition, the probing of stimulated ionospheric irregularities can be obtained from analyses of the signals received at far distance from the HF heater. The paper uses a novel method of scattering of the HF radiation by the heating facility for diagnostics of non-linear effects at the super-long radio paths. Experiments were conducted at three different powerful HF facilities: EISCAT (Norway), HAARP (Alaska), and Arecibo (Puerto Rico) and by using different far spaced receiving sites. The key problems for super-long-range propagation regime is the feeding of ionospheric waveguide. Then the energy needs to exit from the waveguide at a specific location to be detected by the surface-based receiver. During our studies the waveguide feeding was provided by the scattering of HF waves by the artificial ionospheric turbulence (AIT) above the HF heater. An interesting opportunity for the channeling of the HF signals occurs due to the aspect scattering of radio waves by field aligned irregularities (FAI), when the scattering vector is parallel to the Earth surface. Such FAIs geometry takes place over the Arecibo facility. Here FAI are oriented along the geomagnetic field line inclined by 43 degrees. Since the Arecibo HF beam is vertical, the aspect scattered waves will be oriented almost horizontally toward the South. Such geometry provides unique opportunity to channel the radio wave energy into the ionospheric waveguide and excites the whispering gallery modes.

Keywords: artificial ionospheric turbulence, very-long-distance propagation, whisper gallery, ionospheric waveguide, self-scattering

INTRODUCTION

Under the influence of powerful radio emissions on the ionospheric plasma, a variety of non-linear effects occur (Gurevich, 2007). These include electron heating, striction, and thermal parametric instability, stimulated electromagnetic emission, radial drifts in the heated region, and so on. The non-linear interference of the powerful radio wave with the ionosphere produces plasma disturbances and creates the broad spectrum of the inhomogeneities, known as artificial ionospheric turbulence (AIT). The AIT generated by HF heating was first reported by Thorne and Blood (1974) who used the Platteville facility located near Boulder, Colorado, USA. Later studies of AIT were conducted in the next generation of heating facilities such as the European Incoherent Scatter (EISCAT), the High Frequency Active Auroral Research Program (HAARP), Arecibo, and Russian facility Sura. Early diagnostic tools to probe AIT included RF, VHF, and UHF coherent radars and optical detectors. Later they were joined by incoherent UHF incoherent radars at EISCAT and Arecibo, as well as by GPS and LEO satellites, such as French satellite Demeter. Yampolski et al. (1997) conducted experiment at the Sura facility using HF broadcasting station as a probing wave while the UTR-2 Radio telescope (Kharkov, Ukraine) was used as a receiver. In a similar experiment, Ponomarenko et al. (1999) used the heating HF frequency close to multiple electron gyroresonance and found that the spectrum of scattered probe wave experienced strong broadening. It was related to the excitation of small plasma striations. Moreover, stimulated electromagnetic emission (SEE) was successfully used to diagnose AIT (Leyser et al., 1989, 1993, 1994; Carozzi et al., 2002; Thide et al., 2005; Norin et al., 2008, 2009; Bernhardt et al., 2009, 2011; Sergeev et al., 2013; Grach et al., 2016).

AIT leads to the resonance scattering of HF-VHF waves by the irregularities with the scales comparable to the wavelength. In addition, the waves will be focused (or defocused) at scales of the Fresnel zone. Different AIT diagnostics are based on these effects. They include coherent and incoherent radars, vertical ionospheric sounding, and ionospheric radioscapy scintillation technique.

The HF waves radiated by a powerful heater were scattered into the ionospheric waveguide. This waveguide is formed in a valley region between the electron density peaks of the E- and F-layers (Davies, 1989). The waveguide is located at high altitude where the electron collision frequency drops and thus the wave attenuation becomes very low. That allows the radio waves to propagate to long distances. In fact, the waves emitted at HAARP were observed at the Ukrainian Antarctic Akademik Vernadsky Station (UAS) with coordinates 65.25 S, 64.25 W, located many thousand kilometers away from the heating facility.

The key problem of providing the waveguide propagation regime is the feeding of the ionospheric waveguide, which is located in the valley between the E and F ionospheric regions. Then the energy needs to exit from the waveguide at a specific location to be detected by the surface-based receiver. During our studies the waveguide feeding was provided by scattering of the HF waves off the AIT above the HF heater. In the indicatrices

of the resonance signal scattering by AIT a fraction of energy propagates within the sliding angle along the waveguide axes thus providing its feeding.

It is known that effective conditions for the waveguide excitation, as well as signal landing from the waveguide, are created by the regular horizontal gradients which appear during sunset and sunrise in the ionosphere. Such conditions were forecasted for HF heating campaigns by EISCAT and HAARP both located in the North Chemosphere while the receiving UAS site was in Antarctica. **Figure 1** shows the map with the radio-paths EISCAT-UAS, HAARP-UAS, SURA-UAS, and Arecibo-UAS marks the dates when the solar terminator passes through those radio-paths (since the radio link Arecibo-UAS is practically meridional the terminator passes across it during equinoxes). Since HF heating campaigns are irregular, the UAS is monitoring continuously the propagation conditions for HF signals. Since 2002, the systematic monitoring of the probe signals radiated by the two HF time service stations CHU (Ottawa, Canada) and RWM (Moscow, Russia) was conducted. Such observations are useful for the comparable analysis of behavior of signals emitted by the HF heating facilities and for the identification of mechanisms related to the waveguide feeding.

An even more interesting opportunity for channeling of the HF signals occurs due to the aspect scattering of radio waves by the field aligned plasma irregularities (FAI), when the scattering vector is parallel to the Earth surface. Such FAIs geometry takes place over the Arecibo HF facility. Here FAIs are oriented along the geomagnetic field line inclined at about 43°. Since the Arecibo HF beam is vertical, the aspect scattered wave will be oriented almost horizontally toward the South. Such geometry provides unique opportunity to channel the radio wave energy into the ionospheric waveguide and excites the whispering gallery modes (Budden and Martin, 1962; Erukhinov et al., 1975). The whispering gallery modes require only the ionospheric F region curvature, and it does not depend on the E region existence. Those conditions can be fulfilled during the nighttime. If the wave emission can be produced parallel to the Earth surface at the ionospheric altitudes, it can provide the energy to enter and exit from the wave channel. We will show later on in this paper that by choosing the proper conditions of the aspect scattering of the HF signals by FAIs, one can execute such opportunity.

Objectives

This paper is aimed to study the excitation of the ionospheric waveguide due to scattering of the HF heating wave by the artificial ionospheric turbulence (AIT).

Methodology

All three experiments used similar methodology. As a probe wave we used the emission of the powerful HF heater scattered off the AIT. The radiation was observed using coherent HF receivers at very long distance ($\geq 9,000$ km) from the powerful transmitter. In those experiments the heating signals are caused by the resonance scattering of the emission off the decameter scale irregularities being of the order of wavelength of the incident wave. The control of the scattering characteristics is provided by the suitable choice of the heating regime and the Sun illumination of the

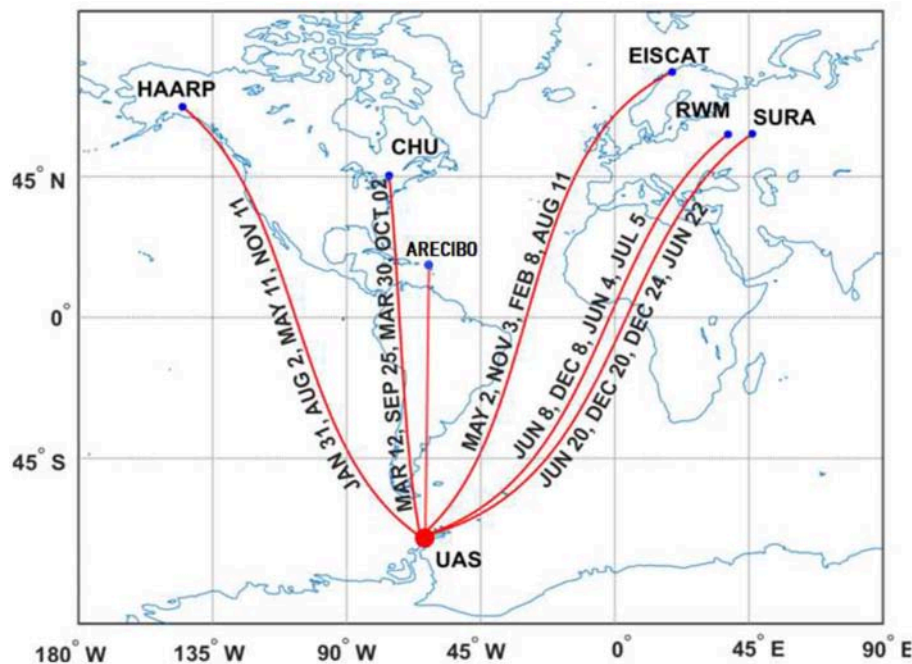


FIGURE 1 | Map of the solar terminator lines and super-long range radio paths.

ionosphere. This way, the ionospheric waveguide can be fed at the beginning of the transmitting line. The energy extraction from the ionospheric wave guide is provided either by refraction off the natural horizontal gradients of the electron concentration (for example during sunset and sunrise) or by scattering off the natural irregularities near the receiver's location. The most convenient conditions for this process occur when the E-layer that screens the energy extraction from the waveguide is absent, i.e., at dusk. Nevertheless, most of the radio link should be Sun illuminated while the E-layer which serves as a lower boundary of the wave guide exists. Consider that the receiver is located in the Antarctic, which is a high latitude region and as such acquires high levels of natural turbulence, it is likely to find here natural irregularities of the decameter scale even during quiet ionospheric conditions.

Experiments were conducted at three different powerful HF facilities: EISCAT, HAARP, and Arecibo and by using different far spaced receiving sites. Note that all discussed results were obtained under quiet geophysical conditions (quiet ionosphere and unperturbed magnetic field). The results of the heating campaigns will be discussed in the chronological order in which they were performed.

SELF-SCATTERING EFFECT DETECTED AT EISCAT

The first successful experiment (Zalizovski et al., 2009) was conducted on 26–30th October 2002 by using the European Incoherent Scatter (EISCAT) facility located near Tromsø,

Norway, with the coordinates 69.35 N, 19.14 E. The heater antenna beam was directed toward the magnetic zenith at 12° zenith angle. The HF operating frequencies varied from 4.0 to 7.95 MHz, the frequencies were chosen to reflect from the F2 peak. Both O- and X-modes were used for the heating. The HF facility radiated power varied from 600 to 900 kW. The ionospheric diagnostic was provided by the EISCAT 929 MHz incoherent scatter radar and ionosonde.

The pump signal was recorded at three spaced receiving sites: at the UAS (Antarctica), at the Radio Astronomy Observatory (RAO, Ukraine), and close to Saint Petersburg (StP), Russia. In addition, the HF radiation of the RWM station of time and frequency service located close to Moscow (Russia) was continuously recorded as a test radiation in Antarctica. The layout of experiment is shown in **Figure 2**.

The experiments were conducted during the fall when sunrise terminator line crossed simultaneously through the interaction region over Tromsø and the most remote receiving site in Antarctica.

We should emphasize the peculiarity of the signals received across all three detection sites. The signal spectrum consists of two components, one narrowband (<0.5 Hz) around the radiated frequency, and the other one broadband (more than 2 Hz) shifted from the radiated frequency by about 2 Hz. It is well-accepted that the narrowband stable signal was formed by the side lobes radiation of the antenna. It propagated along the radio paths by the ordinary hop and multi hops mechanisms. Weak variations of the Doppler shift and amplitude of the narrowband signal were not correlated at different paths. The broadband signal component behaved differently. Variations of the Doppler shift

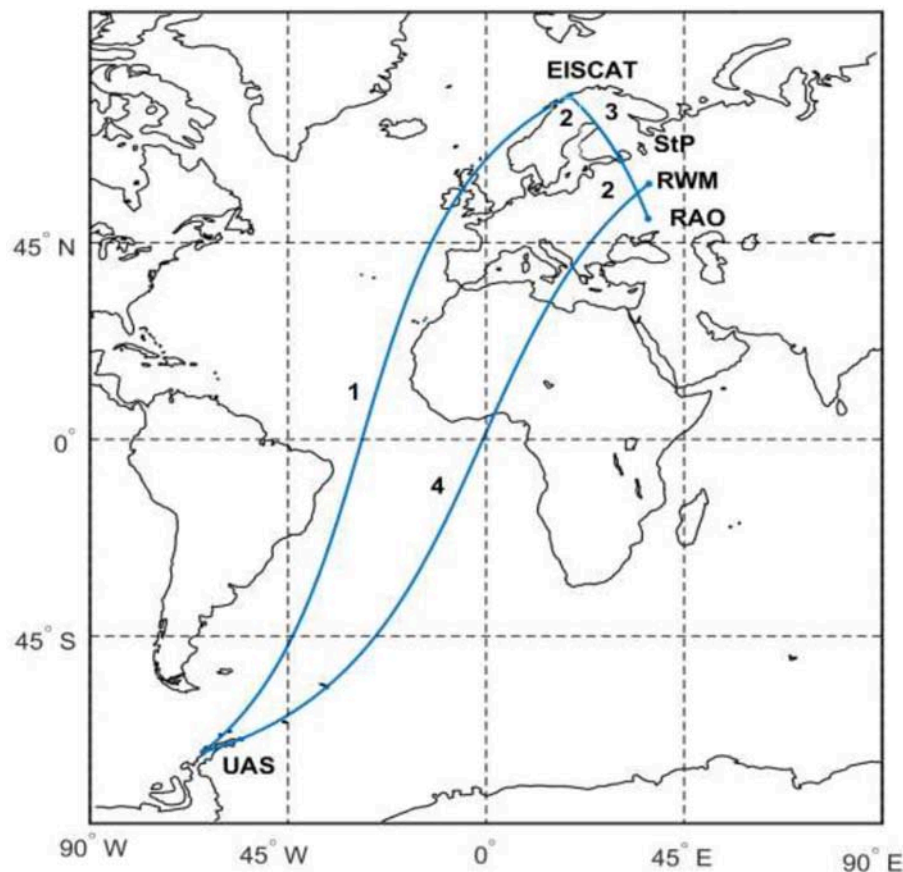


FIGURE 2 | Layout of the experiment. 1: radio path Tromsø–UAS (16,300 km); 2: radio path Tromsø–RAO (2,400 km); 3: radio path Tromsø–StP (1,200 km); 4: radio path RWM–RAO (15,900 km).

and spectral density were well-correlated at all three radio paths. It is illustrated in **Figure 3**, which shows dynamic spectra of scattered components of the signals received at the UAS (a), RAO (b), and StP (c) on 29th October 2002 from 04:34 to 04:37 UT.

EISCAT heating facility radiated two powerful radio waves O mode, which were shifted from each other by 19 Hz. The facility operated in a 5 min on, 5 min off regime. We used two frequencies in order to excite the upper ionosphere plasma oscillations having the combination frequency. The propagation conditions along the radio paths were such that the multiple hop mechanism had not operated. All three panels show well-correlated quasiperiodic variations of the Doppler frequency spectra. The correlation coefficient between the pairs of signals detected by the different receivers was higher than 0.7. Analysis of the geophysical background during the experiment shows that the magnetic field above the HF heating facility experienced similar quasiperiodic variations. Probably they were due to the excitation of the resonance magnetic field micro-pulsations Pc 3. The key question which explains synchronization of the spectra at three different radio paths is how to identify the region which scatters the signals. It is obvious that such region is located in the perturbed area of the ionosphere above the HF heater.

Accordingly, the observed effect was called self-scattering (SS) of the powerful radio wave by the artificial ionospheric turbulence (AIT) (Zalizovski et al., 2009). In the example shown in **Figure 3** quasi-periodic variations of the spectral characteristics of the SS signals were due to propagation of the MHD wave through the scattering region thus causing the AIT modulations. The detected effect of self-scattering was also observed during experiments using Sura heater (Kagan et al., 2006). At the Sura SS was studied using the heating frequency close to the fourth electron gyroharmonic. The HAARP campaign conducted from February 21st to March 3rd, 2008 was very successful by using the opportunities given by the upgraded HAARP facility. In fact, one of the transmitters generated the probe signal having a higher frequency than the ionosphere heating frequency. The probe signal was radiated continuously regardless of the heater operation. The receiving sites were in Antarctica, Alaska, Svalbard island (Norway), Greenland, Ukraine and New England (USA). In some cases, digisondes DPS-4 (Reinisch et al., 2006) were used as receivers. A continuously operated probe transmitter allowed us to estimate the relaxation time of the signal scattering irregularities after the HF heater was switched off (Galushko et al., 2008).

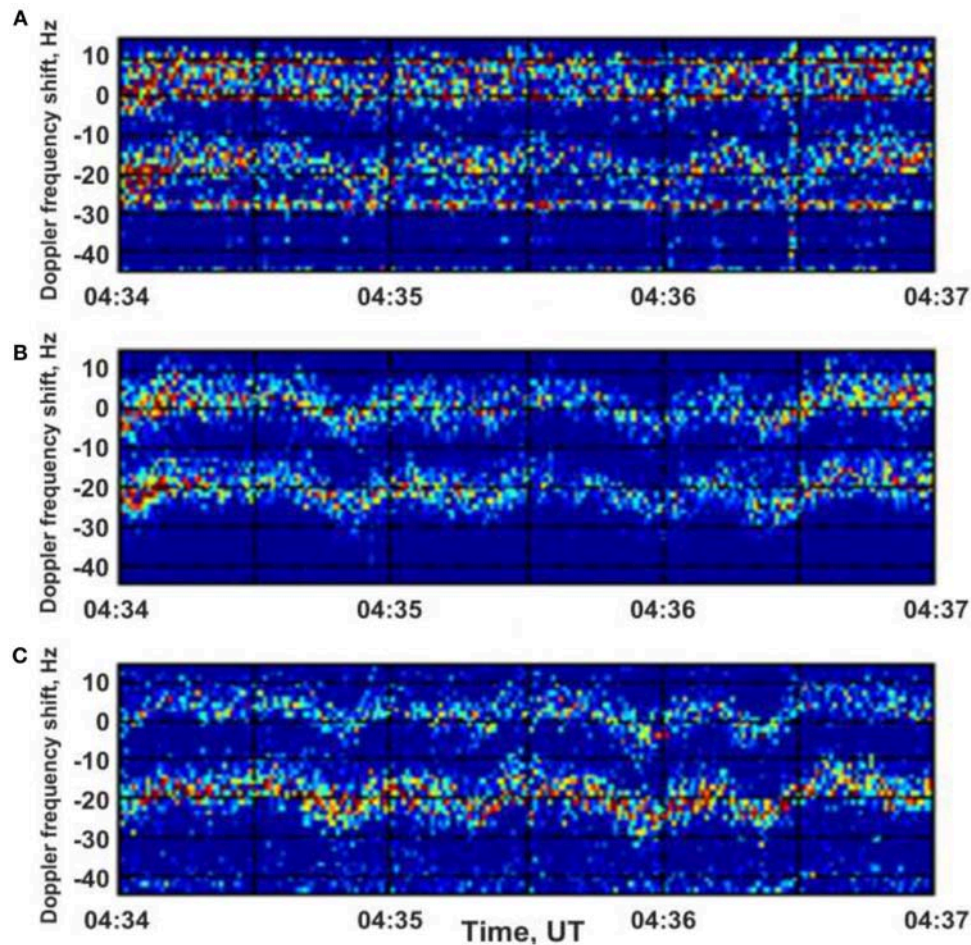


FIGURE 3 | Dynamic spectra for the scattered components of the signal received at the UAS (A), RAO (B), and StP (C) on 29 October 2002 from 04:34 to 04:37 UT.

IONOSPHERIC WAVE GUIDE EXCITATION DETECTED AT HAARP

A successful experiment designated to excite the ionospheric waveguide was conducted over the super-long HAARP-UAS radio path in June 2014. When developing the layout of the experiment, we considered the existing potentialities to control the spatial AIT spectrum. The HAARP heater used O-mode polarization with the HF beam directed at 16° zenith angle and 210° azimuth angle. As described by Najmi et al. (2015) the chosen HF frequency was close to the 4-th electron gyroharmonic (5.6 MHz). It was gradually increased from 5.67 to 5.94 MHz in 30 kHz increments. The heating at each frequency was made by a long pulse of 100 s duration. The pulse consisted of 10 sub-pulses of 10 s each. The ERP was stepped up from 0.7 to 2.5 GW in 0.2 GW increments at each sub-pulse. The artificial ionospheric turbulence excited by the HF beam was probed by the SEE detector located 15 km away from the HAARP site, which was operated by the Naval Research laboratory; and by the HF Kodiak coherent radar located 670 km South West from HAARP. AIT scatters the waves emitted by HAARP into the

ionospheric waveguide. This waveguide is formed between the electron density peaks of the E- and F-regions. The waveguide is located at high altitude where the electron collision frequency drops and thus the wave attenuation becomes low. This allows the radio waves to propagate to super long distances. During the experiment, starting at about 03 UT, i.e., around 7 p.m. local time, the waveguide was oriented along the Earth's terminator. These radio waves were observed on the ground at the UAS 15,600 km away from HAARP. The details of the receiver and of the data acquisition system are presented in Najmi et al. (2015).

The main information regarding the AIT development has been provided by monitoring the SEE. The SEE signals are driven by the non-linear interaction of the injected HF wave with the ionospheric plasma that results in broadband emissions at frequencies different from the injected HF frequencies (Thide et al., 2005). They are usually upshifted or/and downshifted from the heater frequency within a range of 100 kHz. It is known that the BUM is associated with the pumping of 10 cm super small striations (SSS) while the DM is associated with the 7–30 m size striations (Norin et al., 2008). It is illustrated in **Figure 4**, which shows the power spectral densities (PSD) of broadband

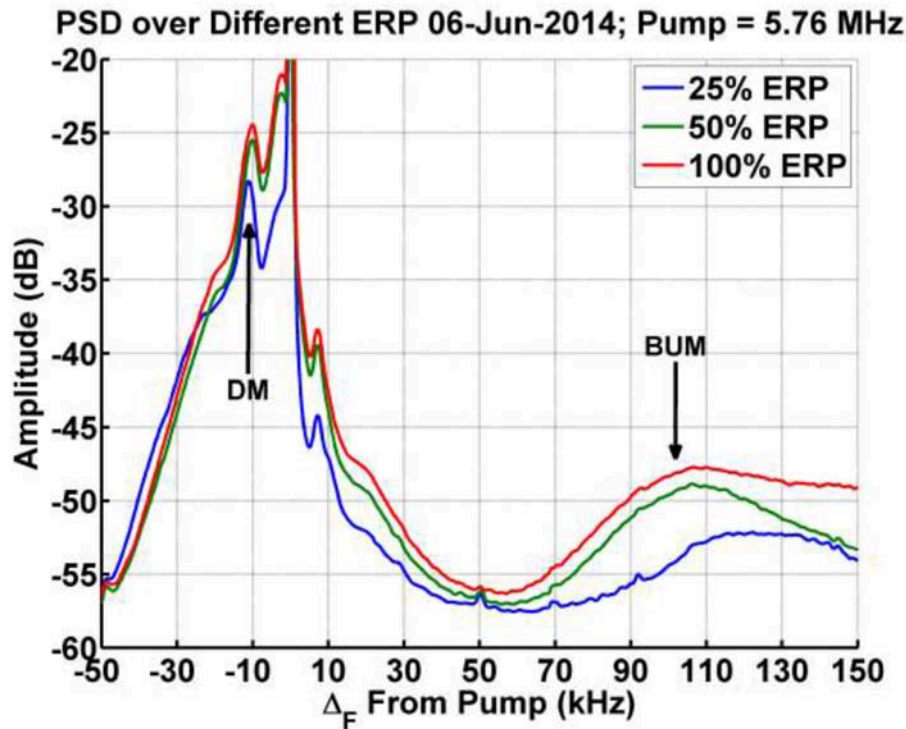


FIGURE 4 | Power spectral densities (PSD) of broadband SEE vs. Δ_F . The traces are averaged over 10 s of the heating time. Variations of the effective radiated power (ERP) are shown in the figure by the color traces. The heating frequency is shown by the highest peak at $\Delta_F = 0$, while the down shifted maximum (DM) and the broad upshifted maximum (BUM) are on the left and right side, respectively [adapted from (Najmi et al., 2015)].

SEE vs. Δ_F . The traces are averaged over 10 s of the heating time. Variations of the effective radiated power (ERP) are revealed in the figure by the color traces. The heating frequency is shown at the highest peak at $\Delta_F = 0$, while the down shifted maximum (DM) and the broad upshifted maximum (BUM) are on the left and right side, respectively.

The development and dynamics of the decimeter AIT irregularities were monitored by the HF Kodiak SUPERDARN coherent radar. The waves radiated by the Kodiak radar have half wavelength ~ 10 m. They are effectively reflected by the artificial striations in the decimeter range. Thus, the radar detected strong scattering due to the ionospheric heating. This is consistent with the velocity of plasma irregularities calculated by the Doppler broadening at UAS, $\Delta f_D = 1.15\text{--}1.40$ Hz. It results in velocity of irregularities $v = \frac{c\Delta f_D}{2f_h} = 30\text{--}35 \frac{\text{m}}{\text{s}}$.

The HAARP HF signals were monitored at the UAS Akademik Vernadsky. **Figure 5** shows the time series of the received power at UAS on June 6th, 2014. Here the heating frequencies are given in MHz. The intensity of the received HF signals vs. elapsed time (i.e., the varying heating frequency) are shown by the color traces. The 8th ($f_h = 5.85$ MHz) and 9th ($f_h = 5.91$ MHz) heating cycles are not shown due to their contamination by interference signals. The blue trace shows the measured data while the red trace is the 10 s moving average. The intensity of detected signal strongly depends on the heating frequency f_h . For f_h slightly above the 4th gyro-frequency, the intensity of the detected signal

was low. The intensity of the detected signal increased with f_h and peaked at 5.79 MHz. We tried without any success to measure scattered HF HAARP signals by ground base detectors in Ukraine and Scandinavia, which emphasizes an important role played by Earth's terminator in the waveguide propagation.

Note that the maximum intensity of the signal received at UAS coincides with the HF heating regime that produces maximum DM SEE. **Figure 6** shows amplitudes of SEE DM and SEE BUM, and the intensity of the HF signals received at UAS vs. the pump frequency. The amplitudes are given in dB normalized by the peak values. In addition, the error bars show the standard deviation of the mean of SNR measured at UAS. The figure reveals that the intensity of the signal detected at UAS is determined by the pump frequency f_h . In fact, when the BUM is the strongest feature in the SEE spectrum, the HF signal at UAS is suppressed. On the other side, when the DM is the strongest feature in the SEE spectrum, the HF signal increases with f_h and peaks at $f_h = 5.79$ MHz. It implies that SEE DM and the amplitude of HF signal detected at UAS peak at the same pump frequency.

We need to emphasize that small scale striations (~ 10 cm), where excitation is associated with SEE BUM, are inefficient scatters of HF waves, which have the half wavelength of about 25 m. At the same time SEE DM shows excitation of the decameter scale striations which efficiently scatter HAARP's HF radiation into the ionospheric waveguide. This radiation is later detected on the ground at UAS.

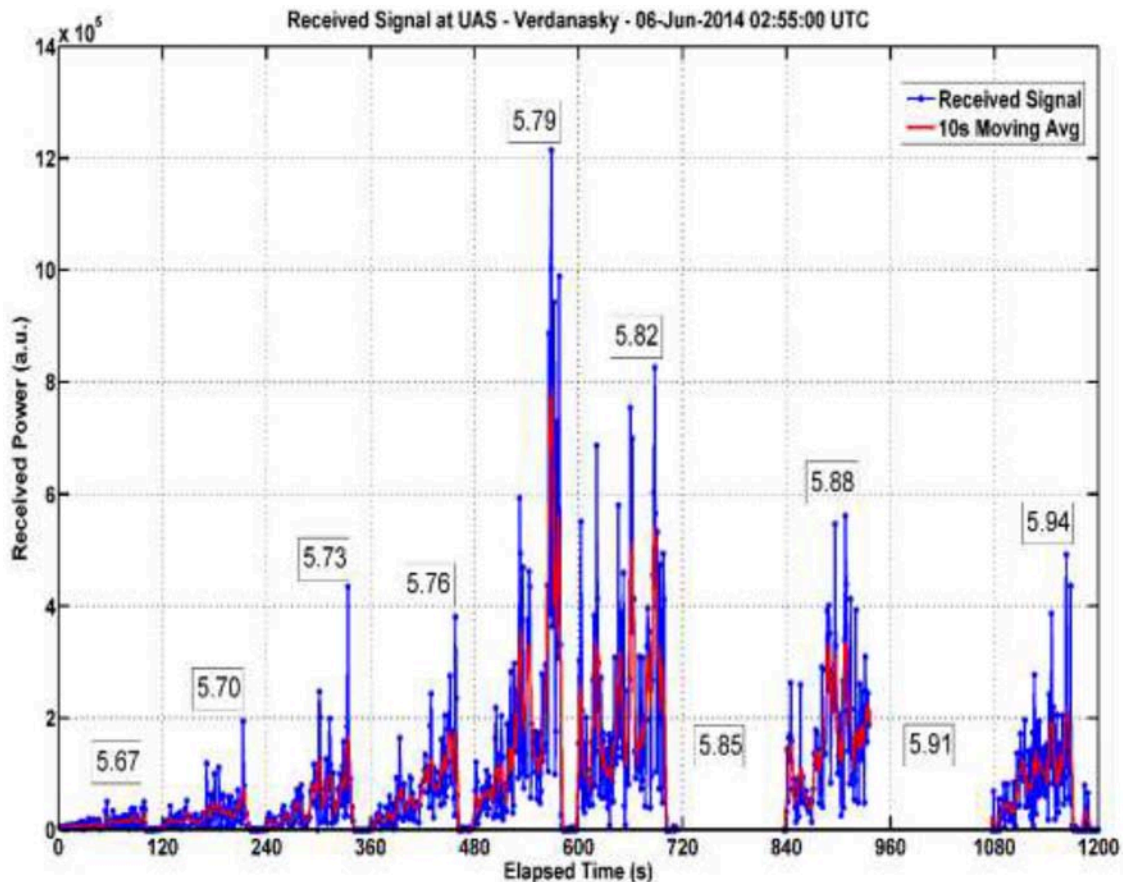


FIGURE 5 | Time series of the received power at UAS on 06/06/2014. Here the heating frequencies are given in MHz. The intensity of the received HF signals vs. elapsed time (i.e., the varying heating frequency) are shown by the color traces [adapted from (Najmi et al., 2015)].

WHISPERING GALLERY EXCITATION DETECTED AT ARECIBO

It is well-accepted that HF heating of the ionosphere generates the FAI. Resonance scattering of the radio waves on FAI is aspect sensitive, and its indicatrix is oriented in the “mirror” direction.

We consider the Arecibo facility as an example. Since the inclination of the geomagnetic field at Arecibo is about 43 degrees and the powerful HF beam is vertical, the aspect scattered wave will be oriented almost horizontally toward the South (see **Figure 7**). Such geometry provides a unique opportunity to channel the radio wave energy into the ionospheric waveguide and excites the whispering gallery modes. Note that in the HAARP–UAS experiment (Najmi et al., 2015) the geomagnetic field geometry has not allowed the generation of whispering gallery modes.

To monitor very-long-distance propagation of the HF signals radiated by the Arecibo facility, we used the receiving facilities at the Akademik Vernadsky station. The respective signal processing enables us to restore Doppler spectra, angles of arrival, time delays and intensities of the signals. This set of parameters allows us to select different propagation modes of the HF signals.

Figure 8 reveals the geometry of the signal propagation on Arecibo–UAS, Arecibo–LFO, and CHU–UAS.

Transmission in the CW mode from the HF time service station CHU (Ottawa, Canada) was used as a reference probe signal for long range propagation (distance between Ottawa and UAS is $\sim 12,300$ km) which allows us to monitor propagation conditions along the meridional direction for daily and seasonal cycles.

The benefits of the Arecibo geometry were mentioned earlier (the aspect scattering wave is horizontal and directed toward South). A shortcoming of the Arecibo experiment is its moderate ERP. The HF facility radiates 600 kW power and has an antenna gain of $G = 22$ dB at $f = 5.1$ MHz thus its ERP is 95 MW. Besides, unlike HAARP, which allows to surf the heating frequency in order to bring it close to proper multiple gyro resonance, which pumps up AIT efficiently, Arecibo operates at two fixed frequencies of 5.1 and 8.175 MHz. Both are far from the multiple gyro resonances. Thus, to describe very-long-distance propagation at the Arecibo experiment, a detailed signal processing is needed.

The experiment took place near the equinox condition on March 17th, 2018 between 22:00 and 24:00 UT (8–10 pm LT). The

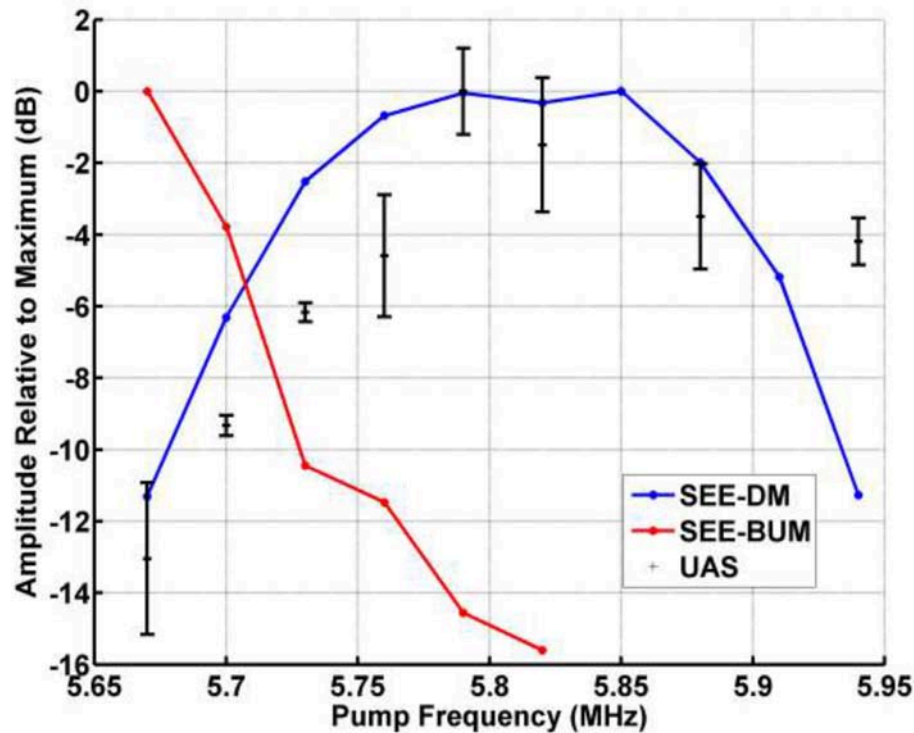


FIGURE 6 | Amplitudes of DM, BUM, and the intensity of the HF signals received at UAS (points with bars) vs. the pump frequency. The amplitudes are given in dB normalized by the peak values. All data points are time averages. Error bars of the HF signals trace are the standard deviation of the mean of SNR measured at UAS [adapted from (Najmi et al., 2015)].

local time at UAS is 1 h forward. The powerful Arecibo HF facility radiated monochromatic signal with the frequency 5,100,067 Hz, O-mode, provided F region heating. We used 20 min heating cycles. During these cycles, for the first 2 min we used -3 dB of the ERP, for the next 2 min we stepped up to -1 dB, for the next 4 min we stepped up to the max ERP (0 dB), while during the next 2 min we stepped down to -3 dB, and at the next 2 min we stepped up to -1 dB, at the next 4 min we stepped up to 0 dB, and finally during the last 4 min of the cycle we stepped down to -10 dB. High frequency stability of the transmitter's generator allows us to conduct coherent spectral processing of the received signals, to select natural as well as artificial ionospheric effects. The frequency of sampling rate was 2 kHz. The main characteristic of the received emission was the amplitude spectra, recovered from 10 s time intervals. The signal intensity was estimated as integral of the square of spectral components in the band $+10$ to -15 Hz with reference to the central frequency $+0.55$ Hz for every 10 s spectrum. Boundaries for the "effective" band were experimentally detected, and outside of these boundaries the spectral density of the ambient noise does not depend on power of the ionosphere heating wave. The carrying frequency 5,100,067 Hz was electronically shifted by $+0.55$ Hz. **Figure 9** shows the spectrogram of the Arecibo signal received at UAS.

Here the narrow band of the signal spectrum corresponds to the carrying frequency. Detailed spectral analysis revealed that the narrow band component practically does not fluctuates with the frequency. Most probably, it forms by the radiation of the

Arecibo antenna side lobes and it propagates along the great circle on Arecibo–UAS route due to the multiple hop mechanism. On the dynamic spectrum on **Figure 9**, the carrying frequency is marked by the continuous blue line. The main signal power ($\sim 97\%$) is concentrated in the spectral band from -5.45 to $+2.05$ Hz, marked by the broken lines. Most probably this part of the spectrum is not related to non-linear ionospheric effects. Intensity "steps" of this component linearly follow changes in the power of the HF heater. The latter was obtained from the experiment log files. Slow increase (~ 17 dB) in the intensity of this signal component over the time of the experiment from 22:00 to 00:05 UT was caused by the changes in the ionospheric conditions along the Arecibo–UAS route. When the experiment passed from daytime to nighttime the radio wave absorption in the lower ionosphere reduced.

It was assumed that the non-linear effects should be noticeable in the broadband component of the spectrum, since this component is formed due to the aspect scattering of the Arecibo radiation by the stimulated field aligned irregularities (FAI).

Spectral power of this component was obtained by extracting the narrow band component at -5.45 to $+2.05$ Hz band from the full power of the received signal in the -15 to $+10$ Hz band. We used the HF heating by 2–4 min pulses, each having constant power. For each of the pulse we averaged the intensity of the received signal in the whole band -15 to $+10$ Hz and in the "broadband." It was assumed that when the Arecibo ERP drops below 25% of its peak, the non-linear effects do not appear. Slow

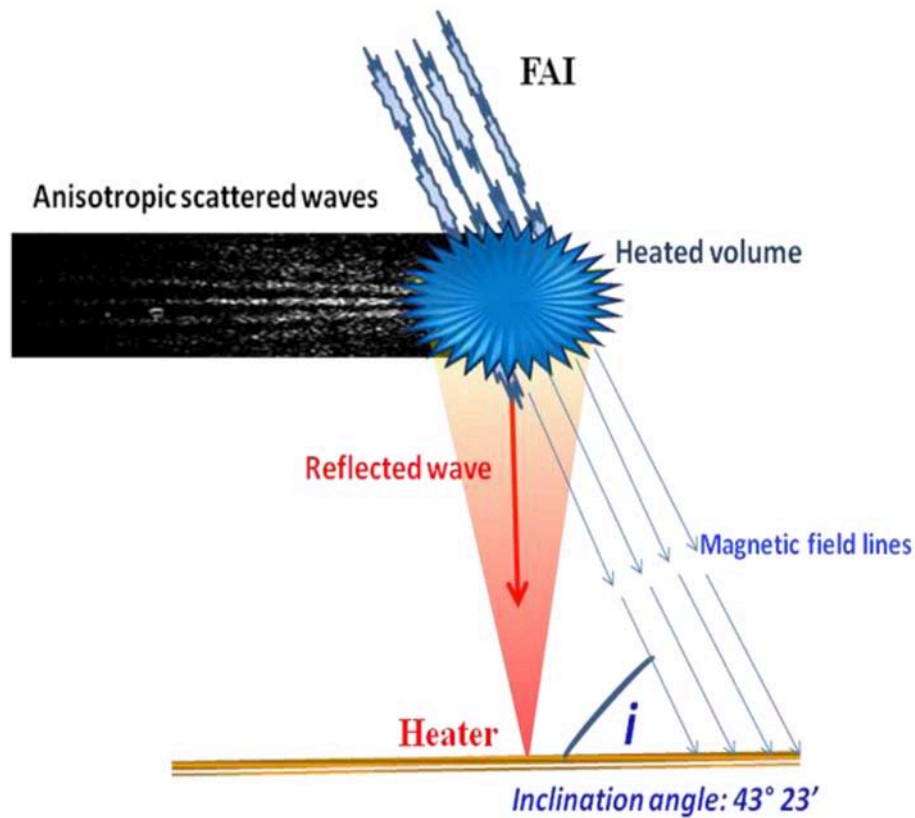


FIGURE 7 | Schematic of the aspect scattering of the radio wave by FAIs above Arecibo heater.

changes of the signal received in the whole band, which were estimated at the minimum level of the heater emission, are related to the variations of the propagation conditions in the ionosphere. Based on this relation we derived a trend, which later was used for detrending and normalization of the HF signals.

The main criteria for the search of non-linear effects was the non-proportional increase of the received signals with the power of the Arecibo transmitter. Averaged values of the intensity of the received signal in the whole band show direct proportionality with the power of the Arecibo HF radiation, when the trend had been excluded. **Figure 10** shows the time series of the power of Arecibo HF radiation, taken from the log file (red histogram); detrended 10 s variations of the full intensity (the black traces). Their mean values for corresponding time interval are shown by the black dashed line; the running second-long variations of the SEE intensity (the turquoise curve); their values averaged over 10 s intervals (the blue curve). Finally, the mean values of the peak SEE power averaged over 1 min pulses are shown by the magenta curves.

Figure 10 reveals that there is no noticeable correlation between the variations of intensity of the Arecibo signal in the whole spectral band received at UAS and the SEE variations. The mean values of the signal level received in Antarctica (the black dashed lines) nearly exactly mimic the power variations of the Arecibo heater according to the log files of the experiment (the red histogram). A similar analysis applied to the narrowband

signals (from -5.45 to $+2.05$ Hz) shows the same results. It allows one to conclude that the main signal components of the radio beam propagating from Arecibo to UAS are most probably excited by the side lobes radiation of the Arecibo antenna. Notice that cross correlation analyses of the variations of the intensity of the main component of the received signal and variations of DM SEE power do not show any noticeable relation between these processes. The broadband component in **Figure 11** behaves differently. First, its intensity does not change proportionally to the log files. Second, under maximum heating power, its intensity variations were well-correlated with the variations of DM SEE power.

Note that the mean intensity steps of the broadband signal are proportional to the heating HF power when the latter is small, although the steps increase non-linearly with ERP when it rises.

Furthermore, we analyzed steps of the SEE intensity under maximum ERP of the Arecibo facility, where we expected that non-linear effects can be developed. Those estimates were checked against the steps of the HF power taken from the log files. There are seven such time intervals during the whole experiment. For four of them the SEE signal steps are by 0.5–2 dB higher than those of the heating power steps. For remaining three cases the SEE signal steps were smaller than the heating power steps. The cross-correlation analysis of the intensities of the broadband and SEE signal revealed the high degree of their resemblance for all four time-intervals (22:27–22:31, 22:39–22:42, 22:59–23:02,

23:19–23:22) when the steps of the SEE signals exceeded those of the heating power. In these cases, the cross-correlation coefficient of the two processes was higher than 0.7. For the remaining three cases the cross-correlation coefficient was <0.4 . Such correlation cannot be accidental. Increase and reduction in the intensity of SEE signals correspond to synchronous changes in the broadband component of the signal received at UAS.

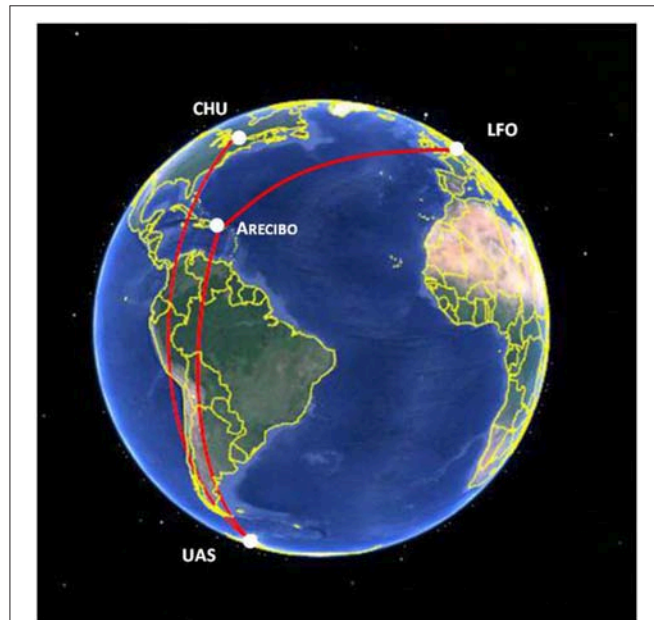


FIGURE 8 | Geometry of the long range radio paths: Arecibo–UAS (9,187 km), as well as CHU (Ottawa, Canada–UAS (12,332 km), and Arecibo–LFO (9,393 km).

Similar spectral and cross correlation analyses were carried out for the signals received at the two test radio paths Arecibo–LFO and CHU–UAS. No relation between variations of different signal components with that of DM SEE power was detected. Therefore, we deduced that the steps of the received HF signal do not depend on spectral frequency band and nearly mimic log files steps; there is no sufficient correlation between variations of the Arecibo ERP and SEE peak power.

Therefore, we can state that conditions of the aspect scattering by FAIs were fulfilled only at Arecibo–UAS radio path. A significant increase (by 5–7 dB) of the level of signal received in Antarctica in phase with the growing intensity of DM SEE is likely to confirm the channeling of energy of the aspect scattered radio wave due to whispering gallery mode. The focusing of the signal was detected by UAS in the nighttime when the E region was absent with no valley. The ionospheric conditions at the receiving site in Antarctica were continuously monitored by the UAS ionosonde, while at Arecibo the incoherent scatter radar was used for this purpose.

CONCLUSIONS

The paper discusses the three experiments conducted over 16 years in which very-long-distance wave propagation was induced by the non-linear effect in the ionosphere irradiated by the powerful HF transmitter.

In the first experiment performed in 2002, the result was serendipitously obtained when using the HF EISCAT facility which radiated monochromatic stationary signals. Three radio paths of different lengths and aligning have detected the broadband spectral components which were strongly intercorrelated. In some case the intensity of these components was either comparable to or higher than the intensity of the narrowband signals formed by the multi-hops of the antenna side lobes emission. The observed effect was caused by self-scattering

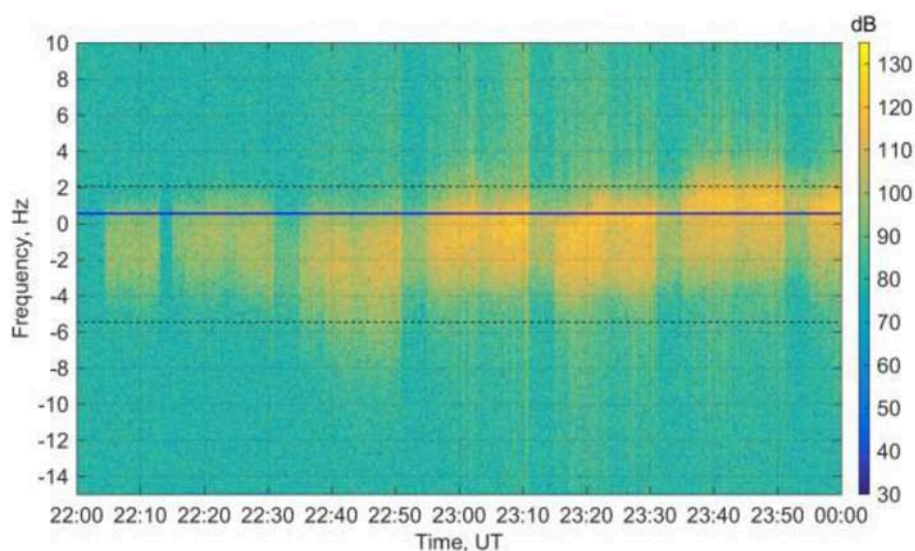


FIGURE 9 | Spectrogram of the Arecibo heater signal received at UAS on March 17, 2018.

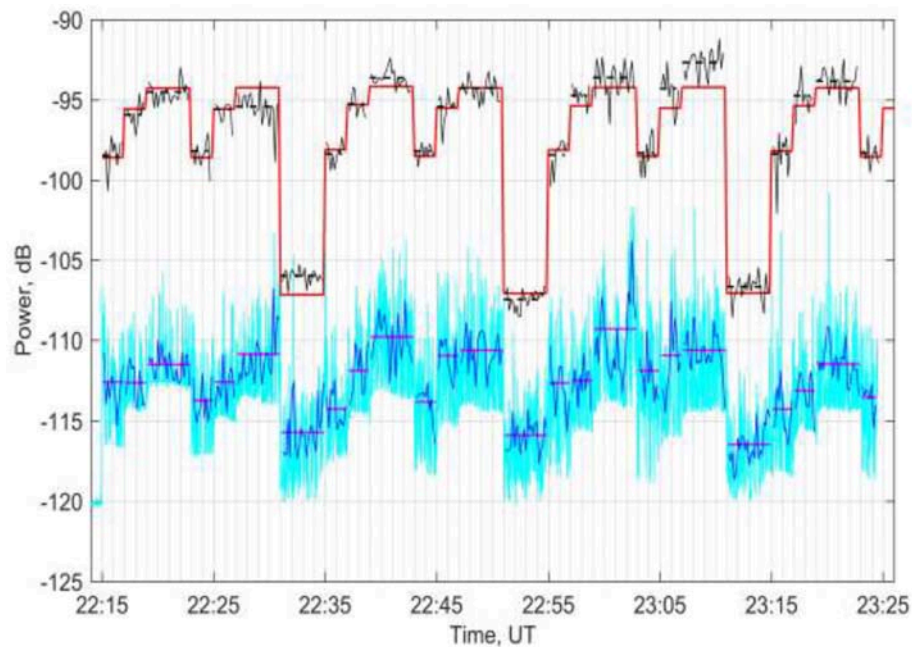


FIGURE 10 | Time series of the Arecibo heater signal received at UAS in the whole spectral band, and the peak intensity of the DM SEE signal received at Arecibo. Red histogram—log file of the heater, black curves—power of received signal at the UAS, black dashed lines an average level of signal power, turquoise curves—intensity of SEE. March 17, 2018.

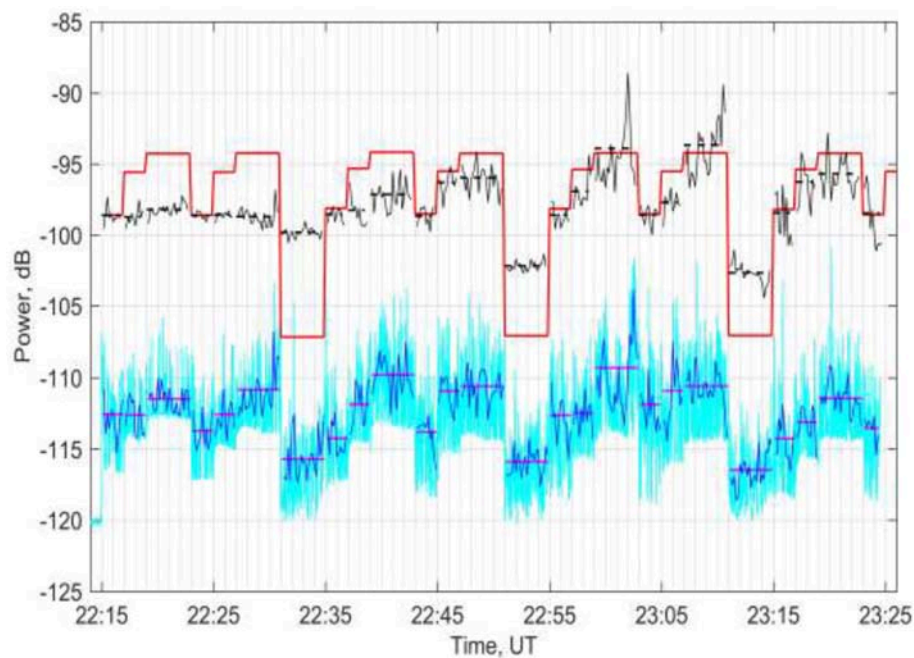


FIGURE 11 | Time series of the intensity of the broadband component of HF signal received at UAS (black curve), noise intensity (the blue curve), red histogram—log file of heater, and turquoise curves—intensity of DM SEE. March 17, 2018.

of the powerful HF emission by the ionospheric irregularities created by the HF emission itself. The broadband spectral component was created by the “secondary” source formed by the

scattering of the Artificial Ionospheric Turbulence excited by the ionospheric heating in the region above the HF facility. Temporal and spatial variations of the AIT region due to the natural

and stimulated drift affect the signal spectra of the scattered component at all the radio links involved. The significant intensity of the self-scattered signal at the very long radio link, which exceeded the intensity of multi-hop narrowband component, allowed us to predict that the ionospheric wave guide can be fed by the secondary source located inside it.

In the second experiment performed in 2014 at the HAARP HF-facility the hypothesis of the artificial feeding of the ionospheric waveguide was proved by using a more sophisticated setup. At the same time, we figured out how to control the AIT efficiency in the domain of meter and decameter scales by sweeping the heating frequency across one of a multiple electron gyroharmonic. Stimulated electromagnetic emission was used to probe the AIT spectrum. During the experiment we detected a sharp increase in the intensity of the HF signal on the very-long-distance radio path HAARP-UAS when the heating frequency approached the 4th electron gyroharmonic.

The third experiment performed in 2018 at the Arecibo facility used the HF energy channeling into the ionospheric waveguide due to the aspect scattering of the radio emission by the field aligned irregularities stretched along the magnetic field lines. The required conditions were fulfilled in the Arecibo heating experiment where the inclination of geomagnetic field is about 43 degrees while the powerful HF beam is vertical, thus the aspect scattered wave will be oriented almost horizontally toward the South. The wave guide was directed toward UAS. Such geometry provides a unique opportunity to channel the radio wave energy into the ionospheric waveguide and excites the whispering gallery modes.

Probing of AIT is obtained from analysis of the radio waves self-scattered into the ionospheric waveguide and then detected at far distance from the heater. However, a probing HF transmitter located in the vicinity of the powerful facility could be very useful for the diagnostic. The frequency of the probing transmitter should be slightly above the heating frequency.

AUTHOR CONTRIBUTIONS

YY: problem formulation and organization of experiments related to detection of super-long-distance propagation of

radio waves generated by HF heaters; GM: organization of the experimental campaigns at HAARP and Arecibo and interpretation of the experimental results; AZ: conducting experiments in Antarctica and making estimates of the radio wave aspect scattering by the natural and artificial irregularities; AK: creating the net of HF receivers at Ukraine, Antarctica, and Arctic operated through the internet and data processing and interpretations; AR: spectral and correlation analysis of the signals propagated along super-long-distances; EN: organizing and conducting experimental campaign at Arecibo; PB: study AIT by using the SEE diagnostic and interpretation of the experimental results; SB: collecting and process SEE data at HAARP and Arecibo; SG: theoretical modeling of generation mechanisms of AIT and SEE and interpretation of the experimental results; AS: processing SEE data at Arecibo; ES: SEE data processing and interpretations.

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Active Experiments Beyond the Earth: Plasma Effects of Sounding Radar Operations in the Ionospheres of Venus, Mars, and the Jovian System

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The operation of powerful satellite- and rocket-born sounding radars is often accompanied by a heating/acceleration of the local electrons and ions. Intense fluxes of sounder accelerated particles were detected in Earth's ionosphere when the frequency of the radar transmitter was close to one of the fundamental plasma resonances: harmonic of the electron-cyclotron frequency, plasma, or upper-hybrid frequencies. Recently it was found that running a sounder in the ionosphere of the non-magnetized Mars results in similar effects. Ion and electron sensors of the ASPERA-3 experiment (Analyzer of Space Plasma and Energetic neutral Atoms) onboard the Mars Express spacecraft discovered acceleration of the local ionospheric ions and electrons from thermal threshold energies to 100's of eV during the active sounding phase of the onboard sounder. ESA and NASA missions being studied or under development to Jupiter (JUICE- JUpiter ICy moon Explorer) in 2022, Europa Clipper in 2023 and to Venus (EnVision) in 2032 and ISRO Venus orbiter in 2023 will also carry powerful sounding radars. The purpose of this study is to investigate what mechanisms can cause acceleration of the plasma particles during operations of the proposed sounding radars in the Jovian system and Venusian ionosphere. Using the results of the previous studies and characteristics of the proposed sounding radars onboard JUICE, Europa Clipper, EnVision, and ISRO Venus Orbiter, we define the optimal conditions for observations of sounder accelerated particles, depending on the local conditions, such as plasma density, composition, and intensity of the magnetic field. The EnVision and ISRO Venus Orbiter radar operations are expected to result in the most pronounced acceleration of ions and electrons, an effect that can be used to improve the local plasma diagnostics.

Keywords: active experiments in space, particle acceleration, ionospheric sounding, Moons of Jupiter, Mars, Venus

1. INTRODUCTION

Ionospheric sounding has been a standard tool for probing the ionosphere for many years. Principles of the sounding are based on the reflection of the radio waves from the ionized component of Earth's upper atmosphere (Appleton, 1927). Ionospheric sounders operate by transmitting a short pulse at a fixed frequency, and then detecting any echoes that are reflected.

These echoes relate to the altitude where the plasma frequency is equal to the wave frequency. By measuring the time delay between the transmission of the pulse and the time that the echo is received, the range to the reflection point can be computed. By sequentially stepping the transmitter frequency after each transmit-receive cycle, the time delay, and hence the range to the reflection point, can be determined as a function of frequency. However, since radio waves return to the ground-based radar from only the portion of the ionosphere below the level of maximum electron density, no information about the ionosphere above this level can be retrieved. To perform topside sounding, ionosondes were first mounted on rockets (Knecht et al., 1961) and later on satellites [Alouette (Lockwood, 1963), ISIS 1, 2 (McAfee, 1969)].

The situation with topside sounding (satellite- and rocket-borne sounders) turned out to be less analogous to bottom-side sounding (ground-based radars) than was expected. Long-term echoes were observed at frequencies of the fundamental plasma resonances: the electron Langmuir frequency f_{pe} , the upper hybrid resonance f_{UH} , the electron-cyclotron f_{ce} and its harmonics nf_{ce} ($n = 2, 3, \dots$) (Calvert and Goe, 1963; Lockwood, 1963). The life-time of these resonances was one or two orders of magnitude longer than the duration of the sounding radio pulse itself. The long durations of the echoes were attributed to the low group velocity of the electrostatic waves that can be generated during the pulse (see Muldrew, 1972 as a review). Moreover, these waves can be reflected by the natural density (for f_{pe} and f_{UH} , McAfee, 1968) or magnetic field (for nH_{ce}) gradients (Fejer and Calvert, 1964). As a result of both these factors, the waves travel a relatively short distance ($10^2 - 10^4$ m) from the sounder with group velocities of $10^3 - 10^5$ m/s and return back, causing long standing echoes.

As it was first pointed out by Oya (1971), it is often the case that the voltage applied to the sounding antenna is comparable to the thermal energy of the electrons. This enables development of a turbulent layer in the vicinity of the sounder, and thus, different non-linear processes should be considered. The first consideration of the stimulated plasma instability and nonlinear phenomena in a framework of the weak turbulence approximation was done by Oya (1971) and further developed by Kiwamoto and Benson (1979) and Benson (1982). Initially, the goal of these studies was to explain the sequence of diffuse plasma resonances, intense echoes observed between harmonics of the electron-cyclotron frequency. The suggested explanation involved cyclotron heating of the plasma surrounding the satellite by the high-power transmitter pulse with subsequent development of the Haris instability (Oya, 1971) or non-linear Landau damping (Kiwamoto and Benson, 1979) that produces observed diffuse echoes.

The first observations that confirmed energization of the plasma by a top-side sounder were made by the Soviet Interkosmos 19 satellite. Galperin et al. (1981) reported observing bursts of superthermal electrons with a mean energy of about 100 eV detected when the onboard high-power radio transmitter was transmitting signals with the tune frequency close to the local plasma frequency. Similar bursts of accelerated electrons and ions were observed also by soft-particle spectrometers (SPS,

Heikkilä et al., 1970) onboard the Canadian satellites ISIS 1 and 2 (James, 1983, 1987). These observations began an intense study of the sounder accelerated particles (SAP) due to the importance of the subject in order to (1) understand the non-linear plasma processes near an active antenna and (2) develop new types of active experiments in space. Different models were proposed to explain SAP observations (see Shuiskaya et al., 1990 for a review). Despite great interest to the physics of SAP, current understanding of the matter is far from being complete, mainly due to the lack of observations.

With greater surprise, a similar phenomenon was recently found in the data collected by particle instruments onboard the Mars Express spacecraft (MEX, Voshchepynets et al., 2018). MEX is the first-ever extraterrestrial planetary mission equipped with both a powerful sounder (for subsurface and ionosphere sounding) and ion and electron sensors (of the ASPERA-3 package, Barabash et al., 2006). The ASPERA-3 plasma measurements cover a period of more than 14 years (more than one solar cycle) and contain multiple observations of SAP. This large set of observations makes it possible to establish statistically reliable dependencies of distribution functions of the accelerated particles on the sounder pulse characteristics (frequency, timing) and the local environmental conditions (plasma density, temperature, composition, magnetic field). The acceleration of electrons and ions by sounders results in the particle beams with known characteristics that can be used for diagnostic purposes. For instance, injection of artificially created beams of energetic electrons is a well-known technique for remote diagnostics of space plasma (Paschmann et al., 2001).

Studying SAP phenomenon is of critical importance to future active experiments in space, a field at the start of the rebirth. The ESA mission to Jupiter – JUICE (JUperiter ICy moon Explorer) – to be launched in 2022 is equipped with both a powerful sounder for subsurface sounding and a comprehensive particle package measuring electrons and ions over a broad energy range. The purpose of this study is to investigate if the sounder onboard JUICE can cause acceleration of the plasma particles in the ionospheres of the large moons Ganymede and Callisto as well as in the Jovian magnetosphere. Other future missions equipped with powerful radars currently known to the authors are Europa Clipper and EnVision.

This paper is organized as follows: The results of observations of the SAP in the ionospheres of Earth and Mars are presented in section 2. Section 3 provides a brief description of the existing theories of the SAP generation. In section 4 we discuss the SAP phenomenon in the context of future space missions.

2. OBSERVATIONS OF SOUNDER ACCELERATED PARTICLES

Overall there were four satellite missions that studied SAP phenomenon in the Earth's ionosphere: International Satellites for Ionospheric Studies (ISIS 1,2, James, 1983, 1987), Interkosmos 19 (Ik-19, Galperin et al., 1981), and Cosmos 1809 (Shuiskaya et al., 1990). The satellites were operational on orbits between 500 and 3,000 km altitudes for different periods from

TABLE 1 | Particle detectors used in SAP studies.

Mission	Instrument	Measured particles	Energy range	Energy resolution	Field of view	Sampling time
EARTH IONOSPHERE						
ISIS 1,2	SPS	Electrons	5.5 eV–13.15 keV	24.7%.	13° × 25°	11 ms
		Ions	5 eV–14.68 keV	35.5%.	13° × 25°	11 ms
IK 19	SF-3	Electrons	10 eV–15 keV	19%.	10° × 30°	20 ms
Cosmos 1809	SF-3M	Electrons	10 eV–10 keV	21%.	20° × 190°	10 ms
		Ions	10 eV–10 keV	28%.	20° × 190°	40 ms
MARS IONOSPHERE						
MEX	ELS	Electrons	1 eV–20 keV	8%.	4° × 360°	28 ms
	IMA	Ions	1 eV–30 keV	7%.	90° × 360°	120 ms

TABLE 2 | Sounders used in SAP studies.

Mission	Length	Frequency range	Pulse duration	Pulse repetition frequency	Power supply
EARTH IONOSPHERE					
ISIS 1	73 m	0.1–20 MHz	86 μ s	45 Hz	400 W
ISIS 2	73 m	0.1–20 MHz	86 μ s	30 Hz	400 W
IK 19	50 m	0.3–15.95 MHz	133 μ s	58.6 Hz	140–300 W
Cosmos 1809	50 m	0.3–15.95 MHz	133 μ s	58.6 Hz	140–300 W
MARS IONOSPHERE					
MEX (MARSIS)	40 m	0.1–5.5 MHz	91.4 μ s	127 Hz	60 W

1969 to 1990. These unique missions were equipped with both powerful radars used for ionospheric sounding and particle detectors that enabled the discovery and studies of the SAP phenomenon. Characteristic of the onboard sounders and particle detectors can be found in **Tables 1, 2**. SAP electrons and ions have typical energy ranges up to 300–500 eV and 100–200 eV, respectively. Electrons show very narrow angular distributions within a restricted pitch angle range around 90°. Ions showed weak pitch angle dependence. SAP are observed in bursts, a burst has a duration of 0.3–3 ms (Galperin et al., 1981). The most intensive electron fluxes were registered when the transmitter frequencies corresponded to one of the resonance frequencies of the surrounding plasma. Accelerated ions are typically detected when the transmitter frequency is near the local plasma frequency. James (1983), based on the measurements of the ISIS 1,2 satellites, showed that accelerated electrons are detected also when the sounding frequency is close to the electron gyro-frequency or its harmonics. Later, Shuiskaya et al. (1990) reported that the SAP electrons can be detected when the sounder operates in the range of the diffuse resonances (a broad area between nf_{ce} when harmonics are below f_{UH}), and between the second and third harmonics of the local plasma frequency.

Very interesting results were obtained from the data collected by the sounding rocket OEDIPUS-C (James et al., 1999; Huang et al., 2001). In this experiment, two suits of instruments were accommodated in two separate platforms connected by an electrically conducting tether that was closely aligned with the Earth's magnetic field during the flight. The forward payload included a high-frequency radio transmitter HEX, which

operated in a frequency range 25 kHz–8 MHz. A synchronized wave receiver called REX was located on the aft payload platform to monitor the characteristics of waves emitted from HEX. Both platforms were equipped with electron detectors that provided measurements of the electron fluxes within the range 10 eV to 20 keV. On the forward platform, the transmission at the electron f_{ce} (and sometimes its harmonics), HEX caused electron acceleration throughout the duration of the rocket flight. The maximal energy of the electrons detected during the OEDIPUS-C experiments was above 10 keV, that is much higher than the SAP energy recorded by any previous mission. The aft payload also detected accelerated electrons in the similar energy range. Simultaneously with the electron receiver on aft payload, strong emissions in a whistler-mode frequency range (200–700 kHz for this experiment) were detected. Unlike the forward payload, the aft payload observed the accelerated electrons with a short delay ($\sim 300 \mu$ s) following the start of transmission and persisted for a short period after the transmitter was turned off.

There are two types of space borne radio sounders. Radars of the first type are used for ionospheric and sub-surface sounding. They are designed to produce distant radio echoes in order to study ionospheric density profiles below the spacecraft or to acquire data about sub-surface structure. Sounders onboard ISIS, Ik-19, and MEX belong to this type of sounders. The second type, known as a relaxation sounder, is a low-power sounder designed to stimulate local plasma resonances. Sounders of this type were flown on numerous magnetospheric satellites such as ISEE 1 (Harvey et al., 1979), GEOS 1 and 2 (Etcheto and Bloch, 1978), CLUSTER (Décréau et al., 2001), and extraterrestrial missions, such as Ulysses (Stone et al., 1992) and Cassini

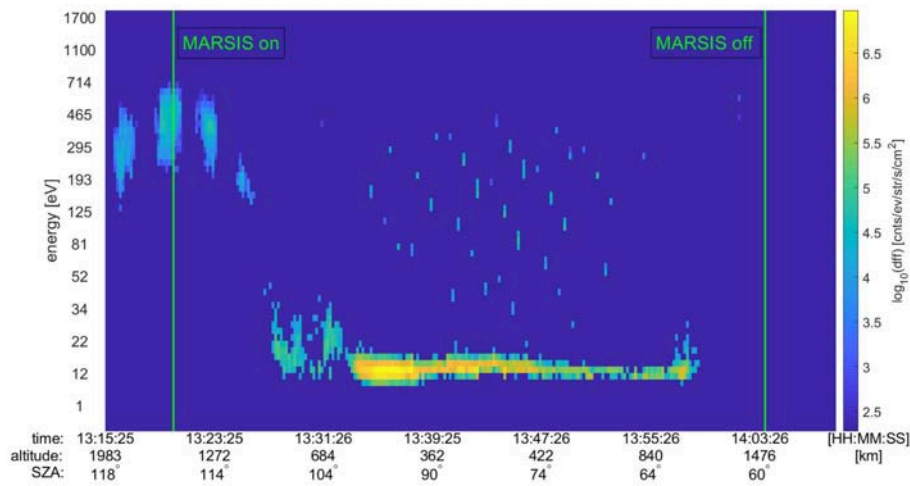


FIGURE 1 | Differential flux of O^+ ions measured by the ASPERA-3 IMA sensors on 22-03-2012 from 13:15 to 14:09 UTC.

(Gurnett et al., 2004). This type of sounders is optimized for measuring densities of very tenuous plasmas (down to 1 cm^3). The working frequency of the relaxation sounders is thus much lower (typically below 100 kHz) than that of the ionospheric and sub-surface sounders. We found no publications on the SAP being detected on the missions equipped with the relaxation sounders. Some of the radars, for instance Radio Plasma Imager (Reinisch et al., 2000) onboard IMAGE satellite, can combine both types of sounding techniques, but IMAGE did not have any electron or ion sensors.

MEX is the first extraterrestrial mission equipped with both a powerful radar (MARSIS - Mars Advanced Radar for Subsurface and Ionosphere Sounding) (Jordan et al., 2009) and comprehensive particle instruments (ASPERA-3, Barabash et al., 2006). This combination of instrumentation enabled studies of the SAP phenomena under plasma conditions other than those in Earth's ionosphere. Characteristics of MARSIS and the ASPERA-3 ion (IMA - Ion Mass Analyzer) and electron (ELS - EElectron Spectrometer) sensors are summarized in **Tables 1, 2**. **Figure 1** shows an example of the sounder accelerated O^+ ions detected by IMA. The shown data were collected during 1 h around pericenter on orbit 10477. MARSIS started operating in ionospheric sounding mode at 13:20:13 UTC and stopped at 14:03:58 UTC. In this mode, the MARSIS operates in the frequency range 100 kHz 5.5 MHz sending 160 $91.7\text{ }\mu\text{s}$ pulses within 1.257 s and then remains idle for 6.285 s. As one can see, a few minutes before crossing the terminator (around 13:39), IMA started detecting intense bursts of energetic ions with energies 40 700 eV on top of low energy 3–5 eV ionospheric plasma background. The bursts were detected throughout the entire MARSIS operation period when the spacecraft was at altitudes below 800 km. The time between two consecutive observations of the accelerated ions is found to be 15, 23, 38, 60 s that coincides with 2, 3, 5, 8 MARSIS repetition times (7.54 s).

The SAP ions are routinely observed when MARSIS operated near pericenter (altitude 250 km) on the day side of Mars.

Preliminary study of the data collected by MEX from 2007 to 2016 showed that 2,528 orbits (of 2,768 available) exhibit signatures of SAP ions. Observations of the accelerated electrons in the Martian ionosphere are much less frequent. Only several hundreds of orbits were found to exhibit signatures of SAP electrons. **Figure 2** shows an example of sounder accelerated electrons detected by ELS together with SAP ions detected by IMA. In the Earth ionosphere, maximum energy of the SAP electrons is several times higher than that of ions, at Mars the situation is different. As one can see in **Figure 2**, the SAP ions are detected with energies higher than 400 eV, while energies of SAP electrons are below 200 eV. Analysis of a large number of observations gives similar results, the maximum energy of SAP ions is twice as high as the energy of SAP electrons (800 and 400 eV, respectively). Ions are detected when the frequency of the MARSIS pulse lies within the frequency range between local plasma frequency and its first harmonic. Electrons are typically detected when the MARSIS operating frequency matches the plasma frequency or one of the harmonics of the plasma frequency (between $2f_{pe}$ and $5f_{pe}$). The sampling times of IMA and ELS are 120 and 32 ms respectively. The sensors cannot resolve individual bursts of the SAPs lasting for 0.1 ms to a few milliseconds. The observed flux increases result from several bursts of SAPs that occur within one sampling period. A MARSIS pulse is $91.4\text{ }\mu\text{s}$ followed by the 7.9 ms sampling time. Therefore, during one IMA sampling time, there will be maximum $120/7.9 = 15$ SAP bursts per a IMA sampling time and $32/7.9 = 4$ per ELS sampling time.

3. THEORIES OF PARTICLE ACCELERATION BY A SOUNDER

A number of explanations of particle acceleration by an active antenna in plasma have been proposed over the last 40 years (see Shuiskaya et al., 1990 and James et al., 1999 for a review). Due

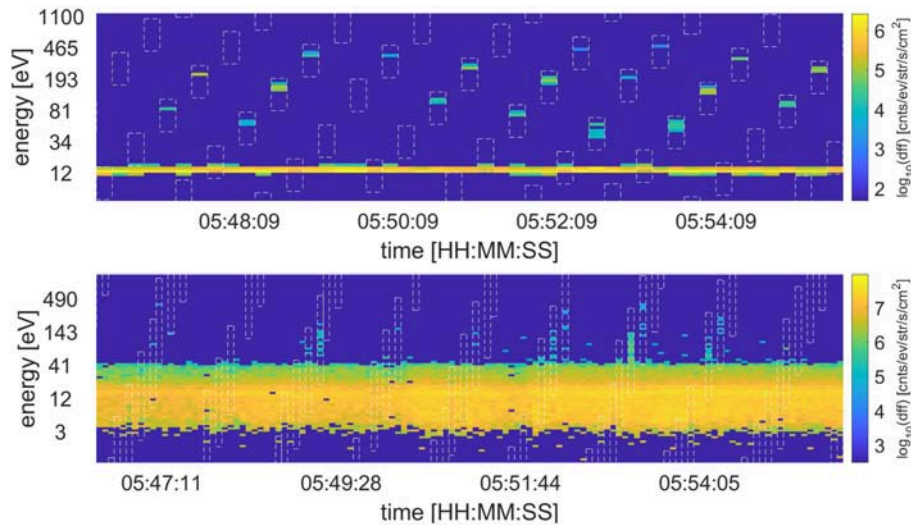


FIGURE 2 | The ASPERA-3 ions (**Top**) and electron (**Bottom**) measurements during MARSIS operations. White boxes indicate the 1.257 s time periods when MARSIS was scanning over frequencies. The fluxes of both accelerated ions and electrons correlate well with the periods when MARSIS was transmitting.

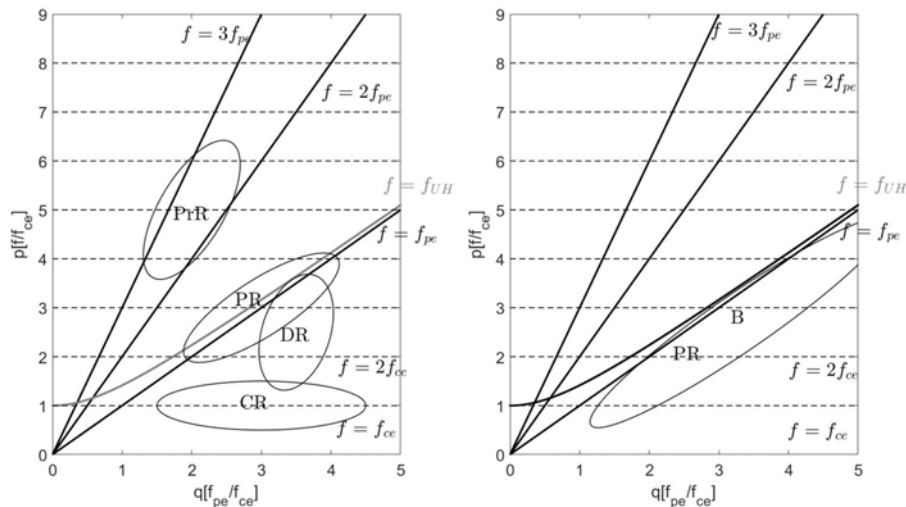


FIGURE 3 | Presentation of the SAP observations by Cosmos 1809 on the $p-q$ diagrams ($p = f/f_{ce}$ and f_{pe}/f_{ce}). The selected areas show the locations of maximum electron (**Left**) and ion (**Right**) intensities together with the indication of the proposed acceleration mechanism. Electron acceleration models: cyclotron-resonance (CR) model (James, 1983), diffuse resonances (DR) model (Benson, 1982), plasma resonance (PR) model (Pulinets and Selegei, 1986), and parametric (PrR) resonance model (Serov et al., 1985). Ion acceleration models: plasma resonance (PR) model (Shuiskaya et al., 1990), wide-band acceleration (B) model (James, 1987). Local electron densities and magnetic fields are summarized in **Table 3**. Adapted from Shuiskaya et al. (1990).

to the fact that SAP are typically detected when the frequency of the transmitter is close to that of the fundamental plasma resonances, several models based on resonance wave-particle and wave-wave interaction have been suggested. Considering only acceleration of electrons, these are: cyclotron-resonance (CR) model (James, 1983), diffuse resonances (DR) model (Benson, 1982), plasma resonance (PR) model (Pulinets and Selegei, 1986), and parametric (PrR) resonance model (Serov et al., 1985). Cyclotron-resonance takes place when $f = f_{ce}$, where f is the frequency of a transmitted pulse. In this model, electron

heating occurs due to Landau damping of electromagnetic oscillations near the transmitting antenna. For the frequencies between f_{ce} and f_{pe} , RF emission can generate electron-cyclotron waves near the antenna. These secondary waves are subject to strong non-linear Landau damping that also results in electron heating (DR model). The PR and PrR models consider electron acceleration in the framework of strong turbulence. In both models, the acceleration occurs as a result of a two-step process. The first stage involves development of the turbulent state near the antenna while the radar is transmitting and

subsequent formation of the cavitons. This can be achieved by the modulation (at $f = f_{pe}$, PR model) or parametric (at $2f_{pe} < f < 3f_{pe}$, PrR model) instability. The collapse of the cavitons after the end of the transmission leads to a burst of Langmuir and ion-sound waves that can effectively accelerate electrons and ions due to the Landau damping. To classify SAP observations in a frequency domain, it is convenient to introduce the parameters $p = f/f_{ce}$ and $q = f_{pe}/f_{ce}$. In this case, each of the proposed mechanisms will occupy a certain area on the $p - q$ diagram, as can be seen in **Figure 3**.

Another effect that is not included in the previously mentioned models is spacecraft charging resulting from the radar operations (James, 1983, 1987). A sinusoidal potential of the radar antenna with respect to the spacecraft during a pulse results in the currents to the spacecraft carrying by either ions or electrons, depending on the potential polarity. If the frequency of the transmitted pulse is lower than f_{pe} , the mobility of the ion and electrons is sufficient to restore the spacecraft potential to an equilibrium value after each half-wave of the pulse. If the frequency of the transmitted pulse is much higher than f_{pe} , neither electrons nor ions are sufficiently mobile to provide currents to change the spacecraft potential. If the frequency of the transmitted pulse is close to f_{pe} , the mobility of electrons is sufficient to charge the spacecraft but the plasma cannot support the sufficient ion current to restore the potential and the spacecraft becomes negatively charged to the amplitude of the pulse. The negatively charged spacecraft attracts and accelerates ions. This mechanism is often referred as wide-band acceleration and it is marked as B in the $p - q$ diagram.

Observations of SAP by MEX on the $p - q$ diagram are shown in **Figure 4**. For the present study, 30 orbits were selected that are characterized by the similar ionospheric environment conditions. Observations of SAP ions correspond to the region of PR and B

mechanisms. Detection of the accelerated ions by IMA is often accompanied by a small decrease in low energy [10–100 eV] electron fluxes as one expects for a negatively charged spacecraft. The induced spacecraft charging model also explains the upper energy limit of SAP ions detected in the Martian ionosphere. As explained earlier, the model suggests that a steady negative DC potential on the order of the voltage applied to the antenna (400 V for MARSIS) is built up on the spacecraft when the transmitter is on. This implies that the instantaneous potential of the antenna with respect to the plasma/spacecraft oscillates from 0 to -800V resulting in the ion acceleration up to 800 eV in agreement with observations. Observations of SAP electrons correspond to the PR and PrR regions on the $p - q$ diagram. It would indicate that the acceleration mechanism could be the resonant wave-particle interaction, but the issue is still under debate. Radar onboard MEX is much less powerful than the radars onboard Interkosmos 19 and ISIS-1,2 spacecrafts (60 W on MEX and 400 W on ISIS-1,2 and 300 W on Interkosmos 19). The case may be that the energy density of the radio waves transmitted by MARSIS is not high enough for the strong turbulence to develop.

4. APPLICATION FOR THE FUTURE MISSIONS

The Jupiter ICy Moons Explorer (JUICE) is an European Space Agency mission that will fly by and observe the Galilean satellites Europa, Ganymede, and Callisto, characterize the Jovian system in a lengthy Jupiter-orbit phase, and ultimately orbit Ganymede for in-depth studies of habitability, evolution, and the local environment (Grasset et al., 2013). It will be equipped with a powerful radar RIME (Radar for Icy Moons Exploration) (Bruzzone et al., 2013). RIME is optimized for the penetration

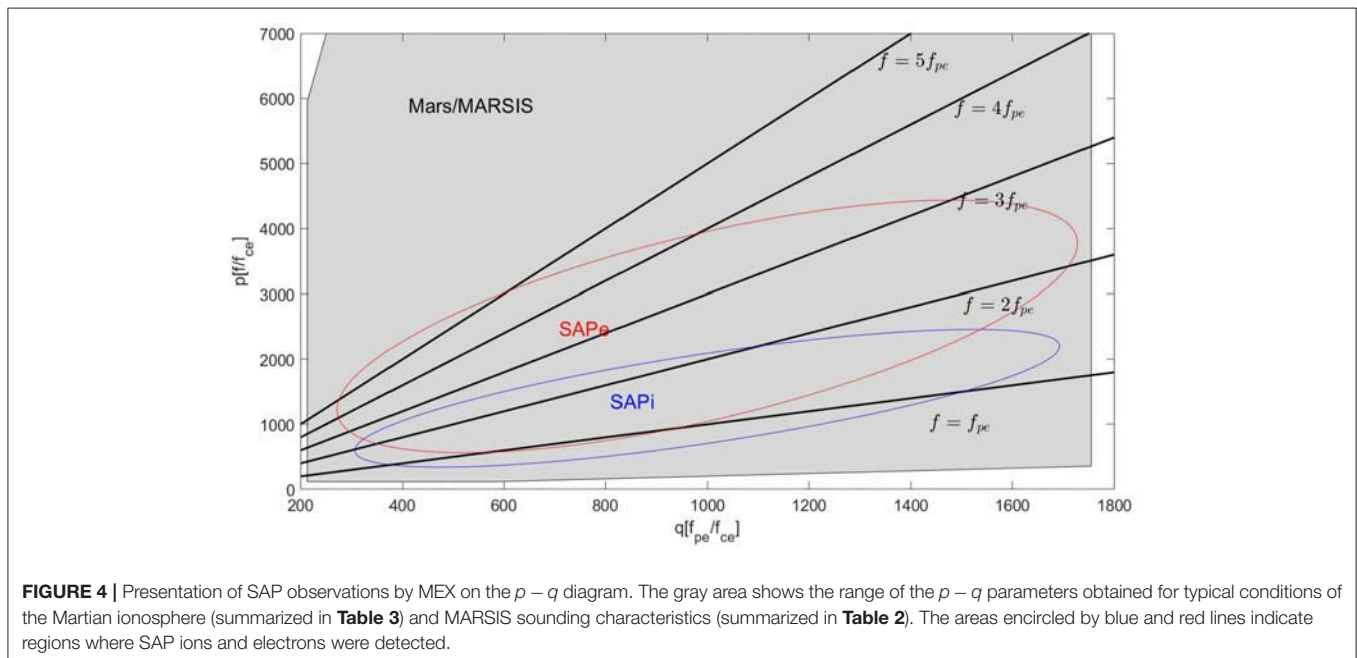
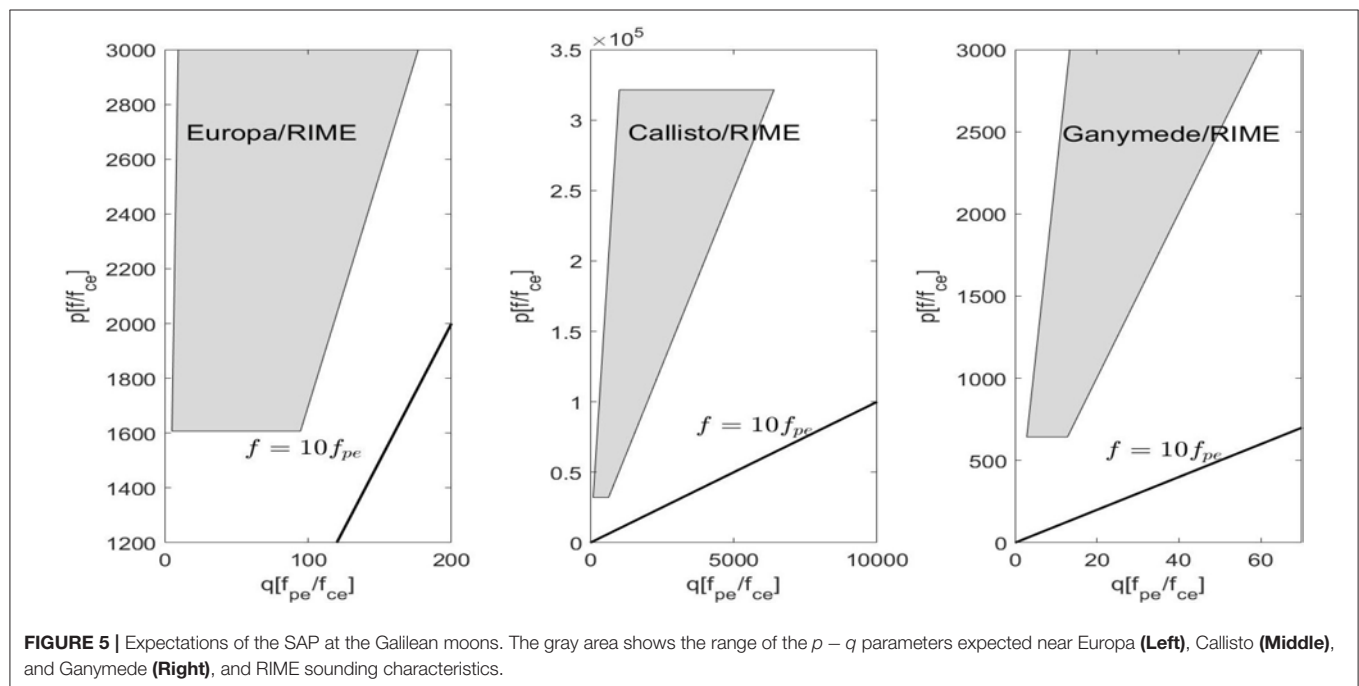


TABLE 3 | Local plasma characteristics (density and magnetic field) typical for SAP observations in the ionospheres of Earth and Mars.

Target	Altitudes	Magnetic field	Electron density	References
IK-19				
Earth	500–1000 km	10–40 μT	$10^4 - 10^5 \text{ cm}^3$	Shuiskaya et al., 1990
MEX				
Mars	300–1000 km	10–30 nT	400 – 3000 cm^3	Fränz et al., 2006
JUICE				
Ganymede	200–1000 km	100–500 nT	20 – 400 cm^3	Kivelson et al., 1997; Eviatar et al., 2001
Europa	400–5000 km	100–200 nT	100 – 3500 cm^3	Kliore et al., 1997; Russell, 2005
Callisto	400–5000 km	1–10 nT	10 – 400 cm^3	Gurnett et al., 2000; Russell, 2005
EUROPA CLIPPER				
Europa	25–5000 km	100–200 nT	$10^2 - 10^4 \text{ cm}^3$	Kliore et al., 1997; Russell, 2005
Europa (plume)	25–5000 km	600–800 nT	500 – 5000 cm^3	Jia et al., 2018
ENVISION				
Venus	259 km	10 nT	$10^4 - 5 \times 10^5 \text{ cm}^3$	Cravens et al., 1997; Donahue and Russell, 1997
ISRO VENUS ORBITER				
Venus	Periapsis at 500 km	10 nT	$10^3 - 10^5 \text{ cm}^3$	Cravens et al., 1997; Donahue and Russell, 1997

Anticipated range of plasma characteristic near Ganymede, Europa, Callisto and Venus for the JUICE, Europa Clipper, EnVision, and ISRO Venus Orbiter missions.



of the Ganymede, Europa and Callisto surfaces up to a depth of 9 km in order to allow the study of the subsurface geology and geophysics of the icy moons (in the search for possible subsurface water). In comparison to ionospheric sounders, radars used for subsurface studies transmit in a higher frequency range. For the RIME, the operating frequency will be set to 9 MHz. Analysis of the plasma conditions near Ganymede, Europa, and Callisto (Table 3) showed that RIME will operate far from any of the plasma resonances. Representation of the RIME characteristics with respect to the local plasma on the q – p diagram is shown in Figure 5. It is highly unlikely that any mechanisms discussed

in the previous section can produce SAP under the conditions expected at the Galilean moons – the RIME frequency is simply too high.

The similar situation is expected for the NASA reconnaissance mission Europa Clipper (Phillips and Pappalardo, 2014). The radar REASON (Radar for Europa Assessment and Sounding: Ocean to Near-surface,) onboard Europa Clipper will have similar characteristics to the RIME and will operate at 9 MHz (Schroeder et al., 2016). Despite the fact that the spacecraft will fly by Europa at much lower altitudes (25 km), the local plasma frequency will still be far below 9 MHz (Figure 6).

This can also be said about the plasma region affected by the water plume rising ~ 200 km above Europa's surface (Roth et al., 2014). The local electron density should be as high as 10^5 cm^{-3} in order to make it possible for SAP to occur.

The most promising missions for SAP studies beyond Earth orbit are EnVision, an orbital mission to Venus proposed to ESA (ESA proposal, Ghail et al., 2018) and ISRO Venus Orbiter proposed by the Indian Space Research Organization (Haider et al., 2018). EnVision will be placed on a circular low altitude orbit (259 km) and will carry a radar designed for subsurface studies (SRS). The radar will work with a central frequency in the range 9–30 MHz for optimal ground penetration capability. Unlike Mars, Venus has a much denser atmosphere. Electron density in the

lower layers of the ionosphere of Venus can reach $10^4 - 10^5 \text{ cm}^{-3}$ (Donahue and Russell, 1997). SRS onboard EnVision will thus operate in a range of $p - q$ parameters that corresponds to the range where MEX detected electrons and ions accelerated by MARSIS (Figure 7). ISRO Venus Orbiter will carry a sounding radar that should have characteristics similar to those of MARSIS on Mars Express. In the ionospheric sounding mode, the radar will be operating in a frequency range between 0.1 and 10 MHz. This frequency range is well suited for SAP generation (Figure 8) near the periapsis (500 km) of the expected ISRO Venus Orbiter orbit. In addition to ion and electron sensors, ISRO Venus Orbiter will carry a plasma wave experiment, Venus Ionospheric Plasma wave detector (VIPER), that will provide measurements of electric and magnetic fields in the vicinity of the spacecraft. The combination of the sounder and particle instruments onboard ISRO Venus Orbiter will enable comprehensive study of the particle acceleration in the vicinity of the active antenna.

5. CONCLUSION

Particle acceleration by sounders in the planetary ionospheres is a type of active experiments of great interest. These phenomena are a result of routine operation of satellite-borne sounders and do not require any additional spacecraft or mission resources, an important factor for planetary missions with scarce resources. On the other hand, it provides additional opportunities for plasma diagnostics (see Voshchepynets et al., 2018). Nominal operations of the radars on the coming missions to Jupiter, JUICE and Europa Clipper, may not result in any notable SAP for typical plasma conditions due to too high radar frequency. However, the comprehensive plasma instrument

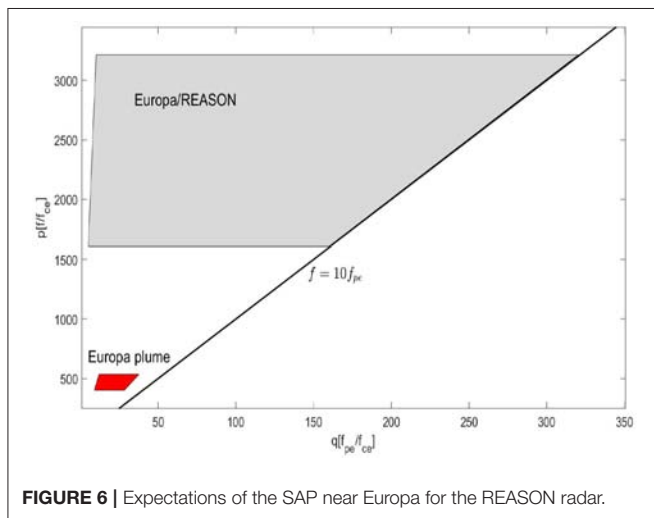


FIGURE 6 | Expectations of the SAP near Europa for the REASON radar.

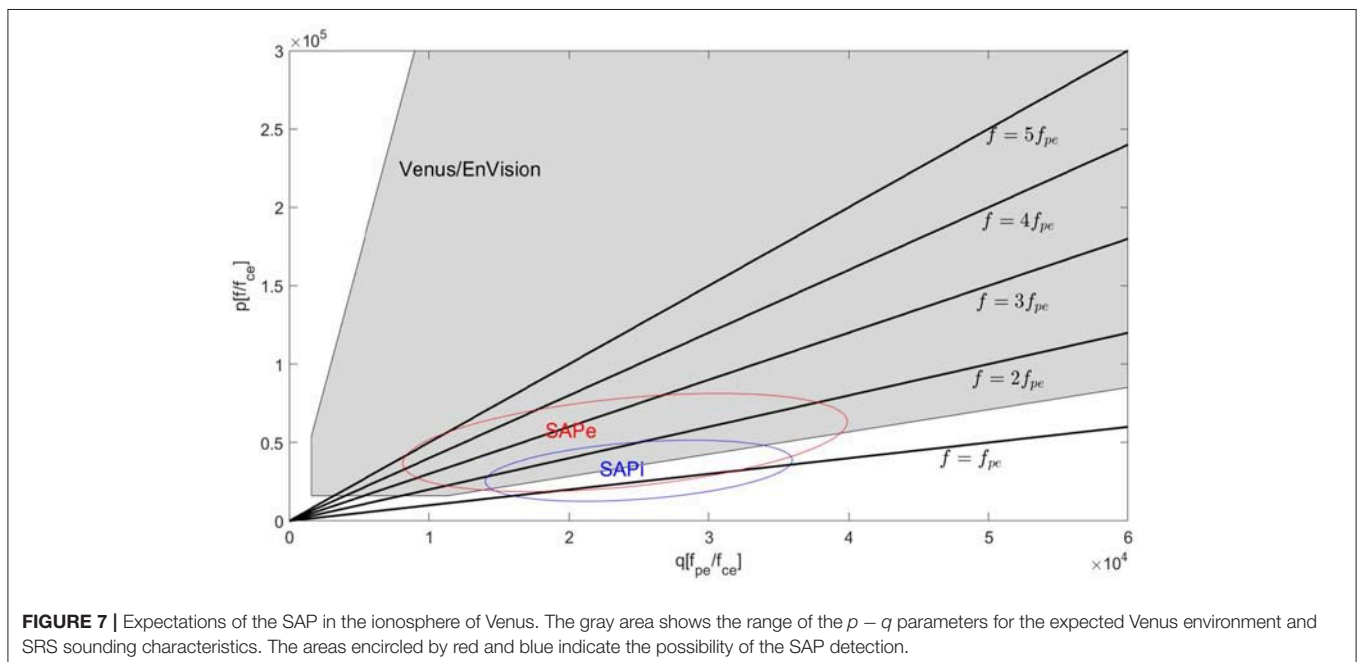


FIGURE 7 | Expectations of the SAP in the ionosphere of Venus. The gray area shows the range of the $p - q$ parameters for the expected Venus environment and SRS sounding characteristics. The areas encircled by red and blue indicate the possibility of the SAP detection.

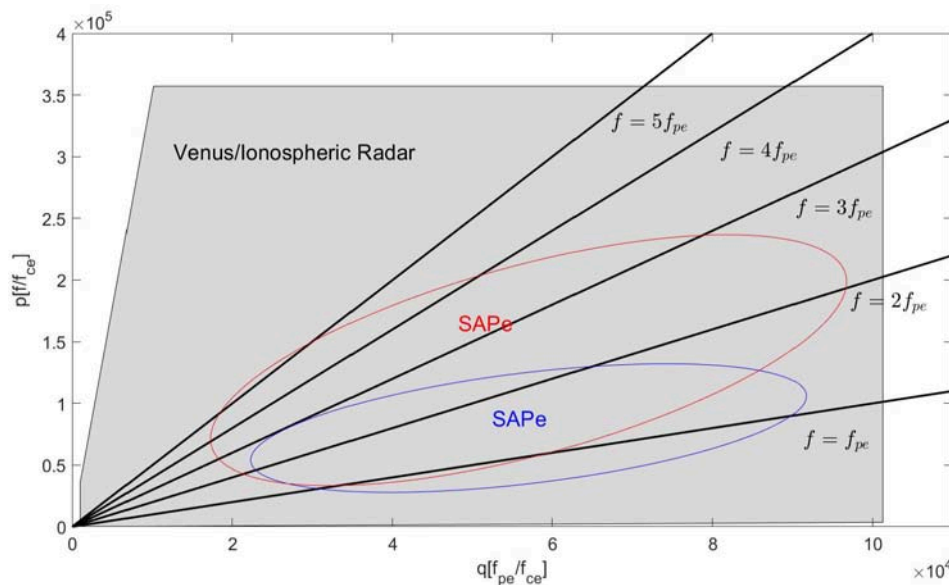


FIGURE 8 | Expectations of the SAP in the ionosphere of Venus. The gray area shows the range of the $p-q$ parameters for the expected Venus environment and SRS sounding characteristics of the radar onboard ISRO Venus Orbiter. The areas encircled by red and blue indicate the possibility of the SAP detection.

PEP (Particle Environment Package, Barabash et al., 2013) onboard JUICE could study plasma modification by an active antenna in the frequency domain far from the main plasma resonances ($f \gg f_{pe}, f_{ce}$).

The sounders on the EnVision and ISRO Venus Orbiter missions to Venus are expected to produce strong SAP fluxes. Currently, EnVision is not equipped with a plasma instrument, but if a conventional ion and/or electron sensor will be added to the payload, they can fully utilize the plasma diagnostic technique made possible by SAP.

Another important aspect of the SAP phenomena is the conclusion on the high negative potential of the spacecraft body. While it occurs for a relatively short time of the order of few ms, it may have effect on the spacecraft platform or operations.

DATA AVAILABILITY

Data used in this study is available on the NASA Planetary Data System and the ESA Planetary System Archive.

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AUTHOR CONTRIBUTIONS

SB and AV contributed conception and design of the study. AV organized the database and performed the statistical analysis. All authors contributed to manuscript revision, read, and approved the submitted version.

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Artificial Aurora Experiments and Application to Natural Aurora

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A review is given of the effects observed during injections of powerful electron beams from sounding rockets into the upper atmosphere. Data come from *in situ* particle and wave measurements near a beam-emitting rocket and ground-based optical, wideband radiowave, and radar observations. The overall data cannot be explained solely by collisional degradation of energetic electrons but require collisionless beam-plasma interactions (BPI) be taken into account. The beam-plasma discharge theory describes the features of the region near a beam-emitting rocket, where the beam-excited plasma waves energize plasma electrons, which then ignite the discharge. The observations far beneath the rocket reveal a double-peak structure of artificial auroral rays, which can be understood in terms of the beam-excited strong Langmuir turbulence being affected by collisions of ionospheric electrons. This leads to the enhanced energization of ionospheric electrons in a narrow layer termed the plasma turbulence layer (PTL), which explains the upper peak. Similar double-peak structures or a sharp upper boundary in rayed auroral arcs have been observed in the auroral ionosphere by optical, radar, and rocket observations, and called Enhanced Aurora. A striking resemblance between Enhanced and Artificial Aurora altitude profiles indicates that they are created by the above BPI process which results in the PTL.

Keywords: active experiments, artificial aurora, electron beam-plasma instability, Langmuir turbulence, Enhanced Aurora

PACS: 94.20.dg, 94.20.Tt, 94.20.wf

1. INTRODUCTION

Aurora, also known as polar or northern lights (aurora borealis) or southern lights (aurora australis), is a natural airglow in the Earth's sky, such as shown in **Figure 1**. As auroras were formerly thought to be the first light of dawn, the name "Aurora" came from the Latin word for "dawn, morning light," while "Borealis" was coined by Galileo in 1619 from the Roman goddess of the dawn and the Greek name for the north wind (Siscoe, 1986). Aurora is seen mainly at high latitudes in the auroral zone, the position of which is controlled by the geomagnetic activity (e.g., Meng et al., 1991). It is produced when fluxes of energetic electrons precipitate along the magnetic field into the upper atmosphere at altitudes below ~ 130 km.

It is thus straightforward to employ electron beams injected from a space vehicle with controlled parameters to explore Artificial Aurora (AA) generated in the upper atmosphere. Such (active) AA experiments have been conducted from sounding rockets and the Spacelab (e.g., Davis et al., 1971, 1980; Hess et al., 1971; Cambou et al., 1975, 1978, 1980; O'Neil et al., 1978a,b; Maehlum et al., 1980a; Jacobsen, 1982; Obayashi et al., 1984; Neubert et al., 1986; Kawashima, 1988; Goerke et al., 1992; Burch et al., 1993). Besides exploring Artificial Aurora, active electron beam experiments, beginning with the Echo



FIGURE 1 | Examples of auroral displays: (a) Corona and (b) rayed arcs (<https://en.wikipedia.org/wiki/Aurora>).

1 experiment (Hendrickson et al., 1971), used injected beams as probes for studying the remote natural processes along the magnetic field. Comprehensive reviews of the “Electron Echo” series and similar experiments have been given by Winckler (1980, 1992) and Kellogg (1992).

The present survey mainly focuses on the AA experiments with primarily downward electron injections. The early experiments were focused on the optical features of Artificial Aurora, then investigation of electron beam-plasma interactions became the main driving force. Because of the limited scope of this review, the setup of different experiments, as well as the electron and ion injectors and diagnostic suites onboard and on the ground, are described only schematically. The presentation is also limited to sounding rocket experiments at altitudes at and below ~ 200 km to avoid the effects of return currents caused by a positive potential of beam-emitting vehicles (e.g., Linson, 1982; Obayashi et al., 1984; Managadze et al., 1988; Frank et al., 1989). The main conclusion of these experiments is that the overall dataset cannot be explained solely by collisional degradation of the beam electrons and requires collisionless beam-plasma interactions (BPI) be taken into account. This stimulated numerous laboratory experiments and theoretical studies discussed during international symposiums on active experiments in space (e.g., Reme, 1980; Grandal, 1982a; Burke, 1983).

This chapter is organized as follows. The salient features of the “classical aurora” limited to collisional impact processes are summarized in section 2. Section 3 gives a review of the effects observed during injections of powerful electron beams from sounding rockets in the ionosphere. Data come from *in situ* measurements of the luminescence, thermal and suprathermal populations, and beam electrons near a beam-emitting space

vehicle, as well as from ground-based optical, radar, and radioemission observations. A brief survey of the BPI theory, which was developed to address the observed unexpected effects, is presented in section 4.

The processes in the near-rocket region are explained by the beam-plasma discharge (BPD) theory (Galeev et al., 1976; Mishin and Ruzhin, 1980a, 1981; Rowland et al., 1981a; Papadopoulos, 1982, 1986; Mishin et al., 1989; Sotnikov et al., 1992) outlined in section 5. The AA rays far from the rocket reveal a special regime of strong Langmuir turbulence in which the wave spectrum in the beam-plasma system, and thus acceleration of suprathermal electrons, is controlled by collisional damping (Izhovkina and Mishin, 1979; Volokitin and Mishin, 1979; Mishin et al., 1981, 1989). The developed theory explained several puzzling features of natural aurora, which is dubbed Enhanced Aurora by Hallinan et al. (1985). Enhanced Aurora (EA) is discussed in the final section. As a rule, only the basic concepts on a semi-qualitative level are given, just sufficient for understanding experimental results. Details and rigorous derivations can be found in the referenced original papers, reviews, and textbooks. Plots and images have been adjusted and sometimes additionally annotated with respect to the originals.

Before presenting the principal experimental results, it is instructive to discuss briefly the “classical” auroral features that follow from collisional interaction of energetic electrons with the neutral atmosphere.

2. BEAM-ATMOSPHERE INTERACTION: “CLASSICAL AURORA”

The collisional or *single-particle* approach (SPA) considers processes of ionization, dissociation, and excitation of atmospheric constituents (N_2 , O, and O_2) by electron impact (e.g., Rees, 1989). In each ionization event precipitating (*primary*) electrons with energies $\varepsilon_b \sim 1\text{--}10$ keV lose energy, $\Delta\varepsilon_{ion} = \varepsilon_{ion} + \varepsilon_s$, where ε_{ion} is the ionization energy and ε_s is the energy of the new-born (*secondary*) electron. Degradation of primary and secondary electrons along the path is usually calculated using Monte Carlo technique. The energy dissipation rate can be estimated from Bethe’s formula

$$\frac{d\varepsilon}{dh} \approx -\frac{\varepsilon}{l_\varepsilon(h)} \propto N(h) \quad (1)$$

Here $l_\varepsilon \approx (\nu/\nu_{ion})\varepsilon/\langle\Delta\varepsilon_{ion}\rangle$ and ν_{ion} is the mean free path and ionization frequency of electrons with the initial energy $\varepsilon = \frac{1}{2}m_e v^2$, respectively; $\langle\Delta\varepsilon_{ion}\rangle \approx 32$ eV is the average energy loss in air; ν is the electron speed; and $N(h)$ is the neutral gas density at a given altitude, h . Equation (1) is valid as long as $\Delta\varepsilon_{ion} \ll \varepsilon$.

The altitude profile of optical emissions at a wavelength λ is determined by the volume emission rate (VER). For the prompt emissions, for which quenching is negligible, the VER reduces to the excitation rate

$$Q_\lambda = 4\pi \int_{\varepsilon_\lambda}^{\infty} \sigma_\lambda(\varepsilon) \Phi(\varepsilon) d\varepsilon \cdot [X] \quad (2)$$

Here ε_λ and σ_λ are the excitation energy and cross-section, respectively, $[X]$ is the density of the neutral species, $\Phi(\varepsilon) = \frac{2\varepsilon}{m_e^2} F(\varepsilon)$ is the differential omnidirectional number flux, m_e is the electron mass, and $F(\varepsilon)$ is the distribution function.

The aurora's color is determined by the wavelengths, λ , of electromagnetic radiation emitted by atoms and molecules in the upper atmosphere, mainly atomic oxygen (O) and molecular nitrogen (N_2), impacted by energetic electrons. The excitation energy is a good indicator of the electron energy distribution. For example, ε_λ for the green- ($\lambda = 557.7$ nm), blue- (427.8), and violet-line (391.4) emissions are about 4.2, 18.9, and 20 eV, respectively. As the excitation cross-sections are small at $\varepsilon > 500$ eV, the auroral glow is mainly determined by suprathermal, $T_e \ll \varepsilon < 500$ eV, electrons with the distribution function $F_s(\varepsilon)$. Here T_e is the electron temperature, which usually does not exceed ~ 0.3 eV.

The suprathermal population is created due to degradation of secondary electrons. In the local approximation, the suprathermal (omnidirectional) flux, $\Phi_s(\varepsilon) = (2\varepsilon/m_e^2) F_s(\varepsilon)$, can be approximated by a power law function

$$\Phi_s(\varepsilon) \approx \frac{3}{2\pi m_e v_c} n_s \left(\frac{\varepsilon_c}{\varepsilon} \right)^{p_s} \quad \text{at } \varepsilon > \varepsilon_c \quad (3)$$

with the spectral index $p_s \approx 3.5$, the density of the secondary electrons $n_s \sim 10 \cdot n_b$ (the density of the primary electron flux), and $\varepsilon_c = m_e v_c^2/2 \approx 6$ eV. At $\varepsilon < \varepsilon_c \approx 6$ eV, the distribution is very sensitive to the neutral composition because of considerable differences between cross-sections for various components. The local approximation is valid until the atmosphere scale height, H_N , greatly exceeds v/v_{il} , where $v_{il}(\varepsilon)$ is the frequency of inelastic collisions. Note that the spectral index for the (directional) magnetic field-aligned flux is $p_s - 1 \approx 2.5$.

Figure 2 shows a typical altitude profile of auroral luminosity calculated for the initial energy $\varepsilon_b = 7.2$ keV using Monte Carlo method (Izhovkina and Mishin, 1979). Apparently, above the peak, at $h > h_m \approx 110$ km, the brightness is proportional to the neutral density N , which is consistent with Equation (1). The peak altitude, h_m , and thickness, Δh_m , can be estimated from the conditions $I_b^{(m)} = I_{\varepsilon_b}(h_m) \approx H_N$ and $\Delta h_m \sim H_N(h_m)$, respectively. Both h_m and Δh_m decrease with increasing ε_b , since the primary electrons penetrate the denser and cooler atmosphere. These features, together with the spectrum $\Phi_s(\varepsilon)$ (3) and the associative VER (2), constitute the “classical aurora” paradigm that is widely used for calculating the auroral E-region conductivity and power released by precipitating electrons.

The next section presents the “anomalous” features of Artificial Aurora that were unforeseen by the collisional approach.

3. ARTIFICIAL AURORA EXPERIMENTS

We start with the observations during the Zarnitza-2 experiment carried out in September 1975 (Cambou et al., 1975; Dokukin et al., 1981; Ivchenko et al., 1981). It is worth noting that Zarnitza 1 and 2 together with the ARAKS (Artificial Radiation and Aurora between Kerguelen and Soviet Union) “North” and

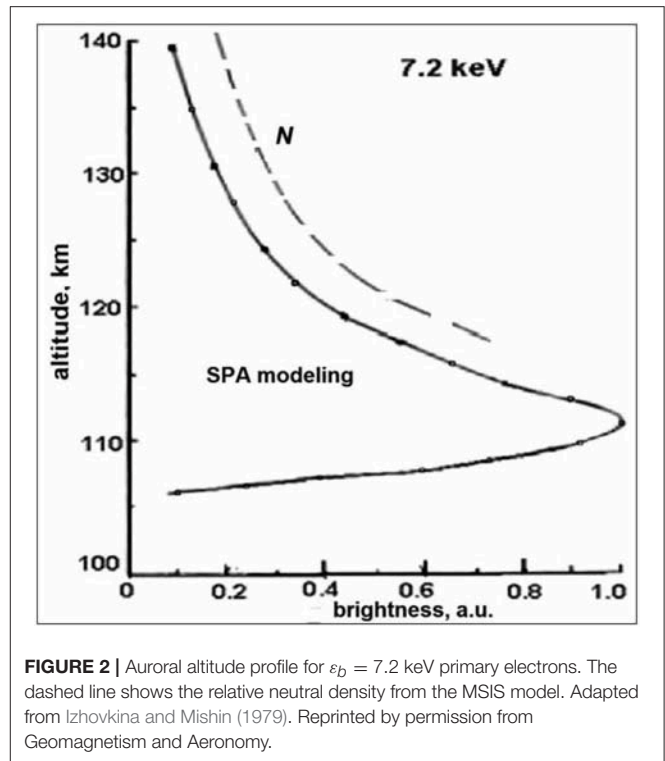


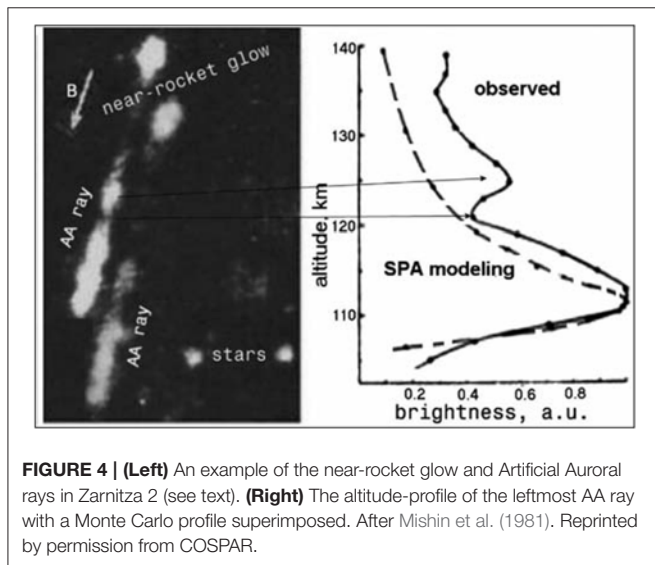
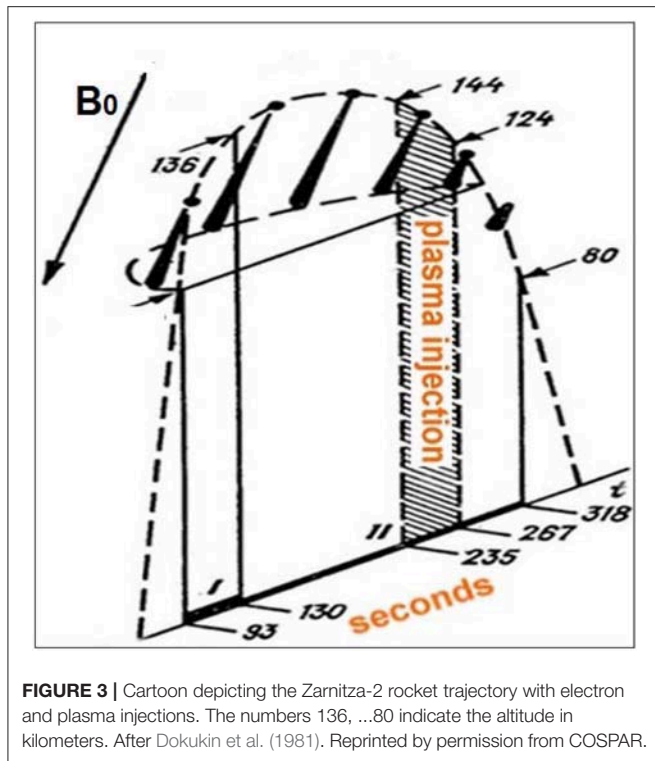
FIGURE 2 | Auroral altitude profile for $\varepsilon_b = 7.2$ keV primary electrons. The dashed line shows the relative neutral density from the MSIS model. Adapted from Izhovkina and Mishin (1979). Reprinted by permission from Geomagnetism and Aeronomy.

“East” experiments constituted the French-Soviet program led by F. Cambou and R. Sagdeev. The electron beam and cesium plasma injectors were put by a rotation-stabilized meteorological rocket into a ballistic trajectory with a 155 km apogee (**Figure 3**). Ground-based diagnostic instruments, including a low-light TV camera, a dual frequency (22.5 and 33.8 MHz) radar, and a broadband (27–51 MHz) VHF radio receiver, provided enough information to describe the effects of injected electrons.

Electron injections were started at $h \approx 109$ km, with the beam energy and current of 9.3 keV and 0.27 A, respectively, and switched to 7 keV and 0.45 A at $h \approx 136$ km. The injection duty cycle comprised of 0.67 and 0.04 s pulses with 100% modulation at 2 kHz and 0.75 s gaps between the pulses. The beam was primarily directed downwards with the initial cone angle of 1° . Because of the rocket's spin with $T_{spin} \approx 3.4$ s, the injection angle, θ_0 , to the magnetic field B_0 varied between the extrema, i.e., $\theta_{min} \leq \theta_0 \leq \theta_{max}$, with θ_{min} (θ_{max}) changing periodically between 40° and 28° (70° and 92°) in 58 s due to the rocket's axis precession.

3.1. Artificial Aurora

In the course of the experiment, a remote low-light TV camera recorded 350–700 nm emissions at 5 frames/s and exposition time of 0.17 s (Ivchenko et al., 1981). All typical auroral lines have been identified in the emission spectrum. Left frame in **Figure 4** is made up of two of about eighty nearly-identical AA images during injections near apogee (Mishin et al., 1981). The luminosity profiles were obtained using microphotometric analysis of the optical emissions along the AA rays and taking the aspect angle condition into account. Each profile features a



bright glow near the beam-emitting rocket and two luminosity peaks far from the rocket. The agreement between the observed and calculated AA lower boundary indicates that the beam energy is close to the nominal (cf. Davis et al., 1971; Rees et al., 1976). That is, the rocket potential is small, which agrees with the measurements onboard (Dokukin et al., 1981).

The AA brightness agrees with Monte Carlo calculations below 120 km (cf. Figure 2), but above the observed luminosity substantially exceeds the collisional level. Besides, the lower peak is slightly wider than the SPA profile, indicating some

additional (~10%) scatter of the beam energy. A slight increase of the energy of some part of the beam electrons indicated by the lowest portion of the profile can possibly be explained by Kainer et al.'s (1972) mechanism.

Now we turn to describe the characteristics of the near-rocket glow (NRG), suprathermal electrons, VHF radioemission, and fast scattering of the beam electrons dubbed “prompt echoes.”

3.2. Near-Rocket Glow and Suprathermal Electrons

Figure 5 shows the AA radiation from the near-rocket glow (NRG) in the course of the Zarnitza 2 and Polar 5 (Maehlum et al., 1980b,c; Grandal, 1982b) experiments, respectively. The variation of the relative radiation flux, F_R/F_0 , from the near-rocket glow (NRG) with altitude and injection angle (Figures 5A–C) does not follow the SPA predictions (Ivchenko et al., 1981). Here $F_0 \approx 2.5 \cdot 10^{18}$ photon/s is the radiation flux from a point source at altitude 100 km, which produces the same flux density, $2 \cdot 10^3$ photon/cm²s, on the ground as the faintest detectable star of the 9th stellar magnitude. On average, the NRG radiation barely changes near apogee. The scatter of the values between 190 and 220 s is due to unstable electric power supply that resulted in the data loss during 220–250 s. Note that fluxes at altitudes ~150 and 115–120 km are of the same order, $F_R/F_0 \sim 10$ –15, while the neutral density changes by more than a factor of ten.

The presence of two almost equal maxima of F_R during one rocket rotation at the extrema of the injection angle, $\theta_{\min} \approx 30^\circ$ and $\theta_{\max} \approx 80^\circ$, also contradicts to the SPA predictions. Taking the average photon energy of 2–3 eV (mainly 391.4-nm photons) gives the radiated power of ~30–45 W. That is, the NRG radiates about one percent of the beam power. These values are of the same order as in the Zarnitza 1 experiment (Cambou et al., 1975) and much larger than the collisional limit. The latter is true for the flux during the whole flight.

Quite similar results have been obtained in the Polar 5 “mother-daughter” rocket experiment (Maehlum et al., 1980b,c; Grandal, 1982b) conducted in February 1976. An electron accelerator on the “daughter” payload produced a ~10 keV electron beam with the maximum current of 0.13 A, which was pulsed at a repetition period 0.4 s. Each pulse comprised five 2-ms sub-pulses separated by 2 ms gaps. The “mother” payload carried a 391.4 nm photometer with the sampling rate of 2.5 kHz and diagnostic instruments for monitoring scattered and secondary electrons, as well as wave effects. The payloads separated slowly, so their distance across B_0 reached about 80 m in the end of flight.

Figure 5D shows the luminescence at 391.4 nm detected by the photometer during beam injections. The observed light level follows the neutral density below 130 km but is fairly constant in the altitude range from 150 km to apogee at 220 km and much larger than that produced by the beam electrons. The latter agrees well with the observations of the electron population in the NRG (Figure 6) with many more suprathermal electrons, the source of airglow, than produced by direct impact.

Indeed, the suprathermal flux at energies less than 1 keV in Figure 6A greatly exceeds the SPA spectrum and has a power

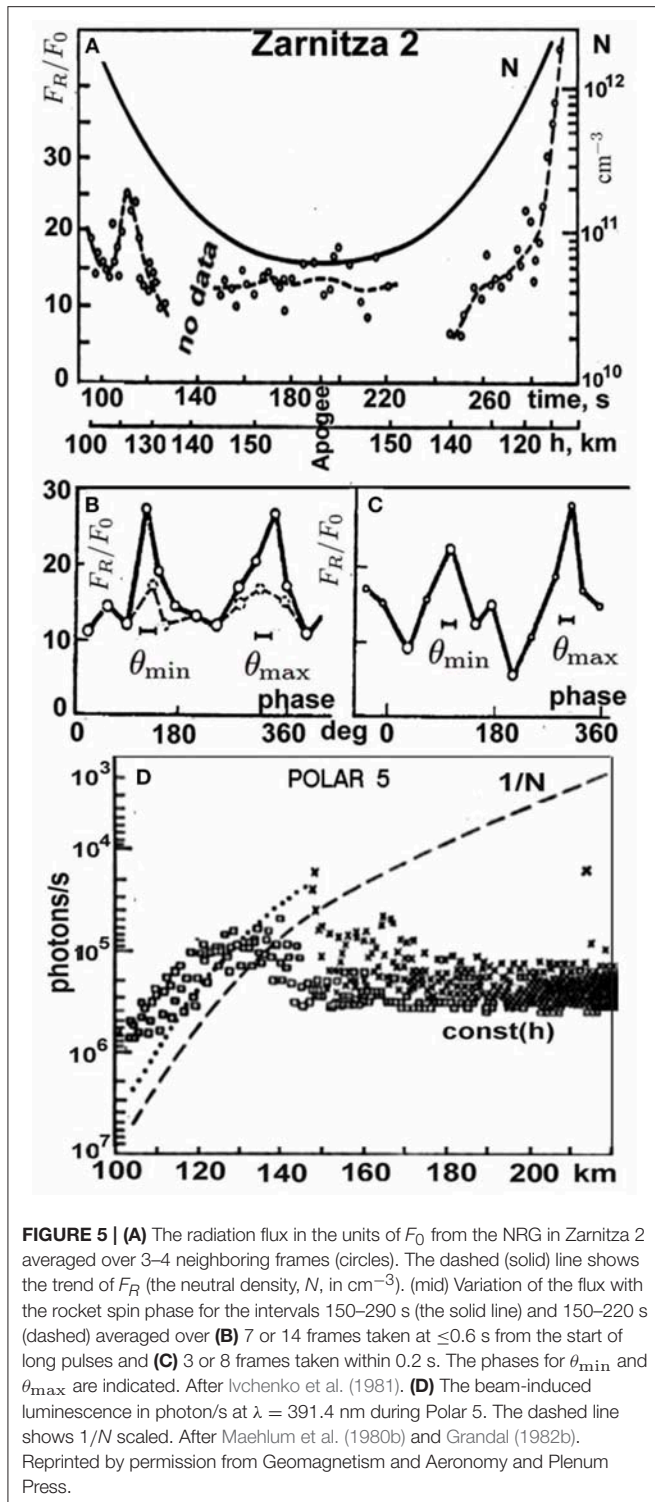


FIGURE 5 | (A) The radiation flux in the units of F_0 from the NRG in Zarnitza 2 averaged over 3–4 neighboring frames (circles). The dashed (solid) line shows the trend of F_R (the neutral density, N , in cm^{-3}). (mid) Variation of the flux with the rocket spin phase for the intervals 150–290 s (the solid line) and 150–220 s (dashed) averaged over (B) 7 or 14 frames taken at ≤ 0.6 s from the start of long pulses and (C) 3 or 8 frames taken within 0.2 s. The phases for $\theta_{\min}$ and $\theta_{\max}$ are indicated. After Ivchenko et al. (1981). (D) The beam-induced luminescence in photon/s at $\lambda = 391.4$ nm during Polar 5. The dashed line shows $1/N$ scaled. After Maehlum et al. (1980b) and Grandal (1982b). Reprinted by permission from Geomagnetism and Aeronomy and Plenum Press.

law spectrum, $\Phi(\varepsilon) \propto \varepsilon^{-1.3}$, i.e., with the spectral index $\sim p_s/3$ (Equation 3). The same is true for suprathermal electron spectra observed during the Echo 5 (Arnoldy et al., 1985; Winckler, 1992) and the Several Compatible Experiments (SCEX) III (Goerke et al., 1992; Bale et al., 1995) experiments in November 1979 and February 1990, respectively. During Echo 5, three electron guns

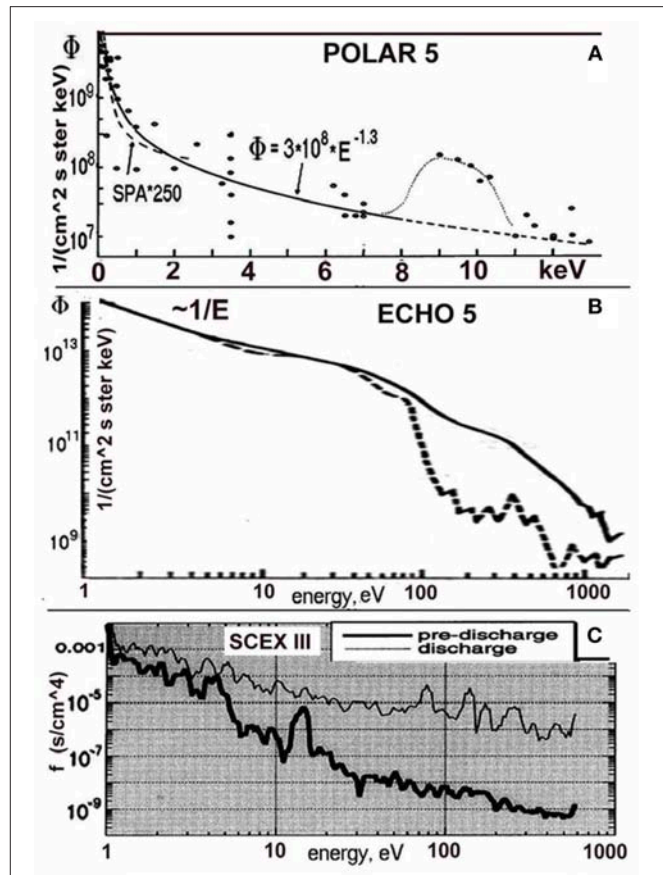


FIGURE 6 | (A) Suprathermal and beam electron spectra in Polar 5. The dashed line shows the SPA spectrum multiplied by 250. After Maehlum et al. (1980b) and Grandal (1982b). Suprathermal electron spectra in (B) Echo 5 with/without (the dashed/solid line) neutral gas injections, and (C) SCEX III. After Arnoldy et al. (1985) and Bale et al. (1995). Reprinted by permission from Elsevier, Plenum Press, and the American Geophysical Union.

injected a 25 keV and 0.6 A electron beam. Each 4-ms injection comprised 0.5 ms turn-on and turn-off pulses and three 1-ms pulses repeated every 20 ms in a 90 pulse “fast” series. Particles were measured every 200 ms aboard the same payload. Due to the electrode voltage sweep failure of the particle detectors at 132 s flight time, only 0.5 and 8 keV particles were measured at 80° and 0° pitch angles, respectively. The suprathermal electron fluxes (Figure 6B) observed earlier in the flight without (the solid line) and with (dashed) cold nitrogen gas injections have the power law spectrum, $\Phi(\varepsilon) \propto \varepsilon^{-1}$, until the flux drops at about 100 eV as the neutral density increased near the rocket. Note beforehand that the latter is in good agreement with the BPI theory prediction (section 4.5).

The SCEX III rocket with a 375 km apogee carried two payloads with a variety of scientific sensors (e.g., Goerke et al., 1992; Bale et al., 1995). An electron gun was located on the aft payload injecting electrons at various energies up to 6 keV and currents from 1 to 60 mA. The neutral density measured on the aft and forward payload showed a significant enhancement over atmospheric models probably caused by severe outgassing.

The plots in **Figure 6C** show suprathermal populations obtained during two 0.16 s sweeps of the retarding potential analyzer (RPA) aboard the aft payload near apogee at 315 s flight time. The “pre-discharge” and “discharge” curve corresponds to the beam current below and above 10 mA, respectively. Above 10–20 mA, the measured (quasi-directional) electron flux averaged over electron energy and 391.4 nm luminosity averaged over 200 s of photometer data (Bale et al., 1995, **Figures 2, 4**) grow nonlinearly with the beam current. The flux’s spectral index is ≈ 2.3 (close to the SPA’s) during the pre-discharge regime and decreases to ≈ 1.2 at greater currents.

The concurrent optical and radar observations in Zarnitza 2 established that the near-rocket glow near apogee had a cylindrical shape with the dimensions of ~ 10 m across and ~ 300 m along $\mathbf{B}_0$ (Dokukin et al., 1981). Similarly, the Polar 5 (Grandal, 1982b), Electron 2 (Jacobsen, 1982), and the U.S./Canadian electron accelerator experiment (Duprat et al., 1983) with mother-daughter payload configurations measured the dimensions of the hot electron/plasma cloud to be ~ 100 m along $\mathbf{B}_0$ and several beam Larmor radii transverse to $\mathbf{B}_0$.

As 391.4 nm emission indicates ionization of nitrogen, the airglow data, consistent with suprathermal electrons at $\varepsilon > \varepsilon_{ion}$, unequivocally point to enhanced ionization taking place in the near-rocket region. This conclusion is supported by concurrent observations of the radar backscatter and intense very high frequency (VHF) radioemission from the near-rocket region at frequencies greatly exceeding the plasma frequency of the ambient plasma (Mishin and Ruzhin, 1980b; Dokukin et al., 1981; Goerke et al., 1992) presented next.

3.3. VHF Radioemission From the Near-Rocket Region in Zarnitza 2 and ARAKS

Figure 7 illustrates the evolution of the VHF spectrum with the injection height and pitch angle during 0.67 s pulses in Zarnitza 2 and the temporal development during 20 ms pulses in ARAKS. In Zarnitza 2, a broadband continuous spectrum of beam-induced VHF electromagnetic waves was detected by a ground-based radio spectrograph with a 27–51 MHz bandwidth (Dokukin et al., 1981). It is seen in **Figures 7A,B** that the radioemission appears during each 0.67 s injection pulse and is modulated by the rocket’s spin. This modulation is due to the variation of the injection pitch angle (cf. **Figures 5B,C**) and also clearly depends on the spin phase. The radio burst near 50 MHz of a 5–10 MHz width at the beginning of the injection pulse is a rapid drift from lower to higher frequencies within the spectrograph sweeping time of 20 ms.

During the first injection regime, the maximum frequency averaged over several long pulses between 120 and ≈ 130 km decreased with altitude as $N^{1/4}$. After transition to the second regime and until the plasma generator was turned on, f_{max} is approximated as follows

$$f_{max} = \text{const} \left(\frac{I_b}{v_b} \frac{N}{V_{R\perp}} \right)^{1/2} \quad (4)$$

Here $V_{R\perp}$ is the rocket’s speed across the magnetic field, which differs significantly between the upleg and downleg parts of the

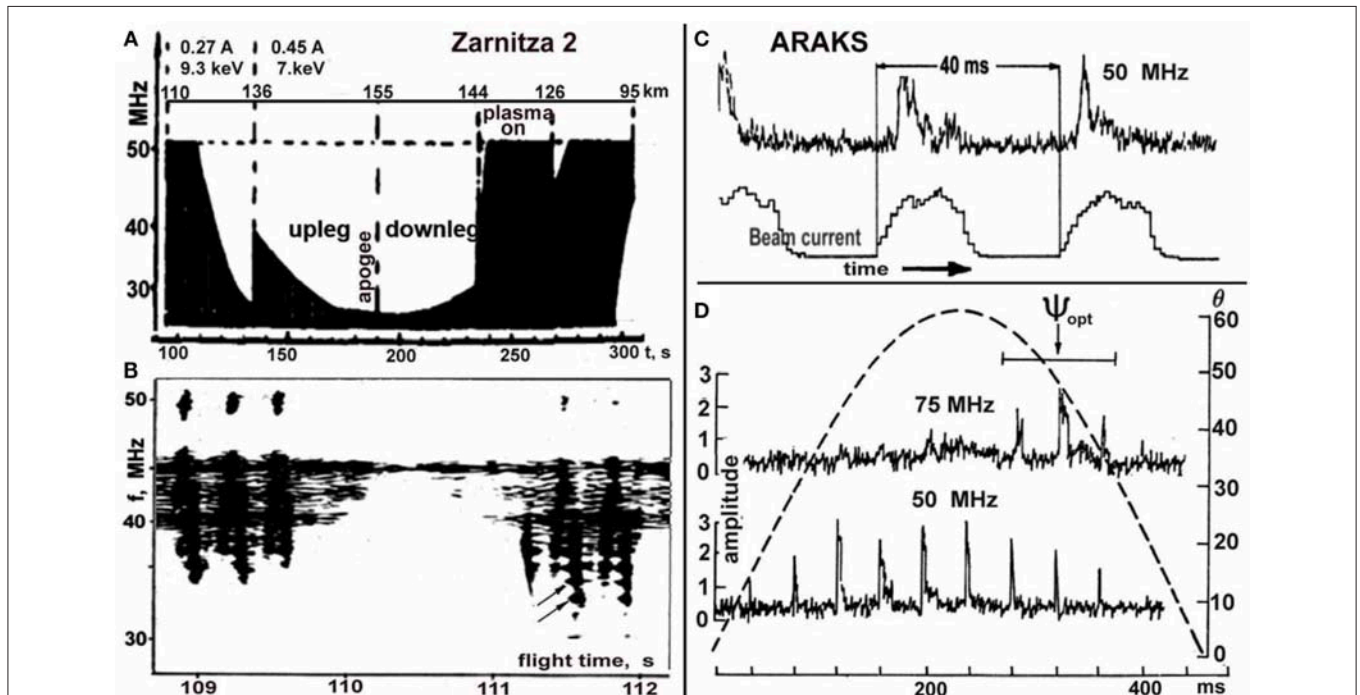


FIGURE 7 | VHF radioemission in Zarnitza 2 and ARAKS: Variation of the broadband spectrum in Zarnitza 2 with **(A)** altitude and **(B)** the rocket’s spin (one long pulse per 2 revolutions). Arrows indicate the spectral maxima at $(s + 1/2)/f_{ce}$. Adapted from Dokukin et al. (1981). **(C)** Delay of radioemission at 50 MHz relative to the start of 20 ms pulses in ARAKS and **(D)** dependence of radioemission at 75 MHz on the phase of the rocket’s rotation. The dashed curve shows the beam pitch angle θ_0 . After Mishin and Ruzhin (1980b). Reprinted by permission from COSPAR.

trajectory. The dependence of the NRG plasma growth on $V_{R\perp}$ was predicted theoretically by Galeev et al. (1976).

Assuming that $f_{\max}$ (4) is the increased plasma (Langmuir) frequency, $f_{pe} \approx 9\sqrt{n_e}$ kHz, gives the plasma density in the NRG, $n_e \propto I_b N / v_b V_{R\perp}$. The spectral density of the electromagnetic flux was at least 10^{-20} W/m²Hz, which exceeds more than ten times the detection threshold. Assuming a point source at about 300 km distance gives ~ 100 W of the total radiated power for a 10 MHz band emission, i.e., ~ 3 percent of the beam power, which is of the same order as the radiated optical power. As shown in **Figure 7B**, the spectrum has maxima (buddges) centered at half-integral harmonics of the electron gyrofrequency, f_{ce} , i.e., $f_s \approx (s + 1/2)f_{ce}$, with integers $s \geq 20$ to ≈ 30 (Dokukin et al., 1981, **Figure 10**). This “gyro structure” is clearly seen only in a certain height range from 110 to 128 km and 105–115 km in the first and second injection regime, respectively. Near apogee, the radioemission at ≥ 27 MHz appears only twice during one rocket's rotation as a brief pulse. At the same time, the VHF radar data show radioemission at the radar frequency of 22.5 MHz during the whole injection pulse, i.e., two rotations of the rocket.

The temporal development of radioemission and its dependence on the phase of the rocket's spin was studied in details during the ARAKS “East” experiment using five narrowband receivers located at the launching site (Mishin and Ruzhin, 1980b). A 0.5 A and 27 or 13 keV electron beam was injected at various angles, φ , with respect to the axis of the rocket. As in Zarnitza 2, the spectral density of the VHF electromagnetic flux was at least 10^{-20} W/m²Hz. **Figures 7C,D** illustrates the features of 50 and 75 MHz emissions generated by 20 ms injections of 13 keV electrons at $\varphi = 30^\circ$ (pitch angles $\theta \approx 0^\circ$ – 60°) in the altitude range of 111.2–103.7 km.

One can see the time delay of a few milliseconds with respect to the start of 20 ms pulses. The time delay for 75 MHz emissions is greater than that for 50 MHz by a few ms (Mishin and Ruzhin, 1980b, **Figure 6**). The generation of the higher-frequency emission depends on the phase of the rocket's rotation, ψ , which is counted off from the axis aligned with $V_{R\perp}$. This dependence results from rotation of the beam guiding center with the spinning rocket, $\Delta V_{\perp} \sim 2\pi r_{cb}/T_{spin}$, where $r_{cb} = v_b \sin \theta / \omega_{ce}$ is the beam Larmor radius (Mishin and Ruzhin, 1980a). The total speed of the guiding center across B_0 amounts to $V_{\perp} = V_{R\perp} + \Delta V_{\perp}(\psi)$. It is minimized at the “optimum” phase, $\psi \rightarrow \psi_{opt}$, when the guiding center moves against $V_{R\perp}$. In the ARAKS experiment, $\Delta V_{\perp}(\psi)/V_{R\perp}$ reached up to 30%.

3.4. Prompt Electron Echoes

Besides the enhanced number of suprathermal electrons in the NRG, strong scattering of beam electrons occurs in the vicinity of a beam-emitting payload (Hendrickson et al., 1971; Winckler et al., 1975; Gringauz et al., 1980; Lyakhov and Managadze, 1980; Maehlum et al., 1980c; Winckler, 1980, 1992; Arnoldy et al., 1985; Wilhelm et al., 1985). This effect is dubbed prompt electron echoes (PEE) as the backscattered electrons are detected within ≤ 100 ms even during upward beam injections. **Figure 8** shows the results from the Echo 5 and ARAKS experiments (Gringauz et al., 1980; Arnoldy et al., 1985). Clearly, above 140 km the scattered flux is almost independent of the neutral density, while

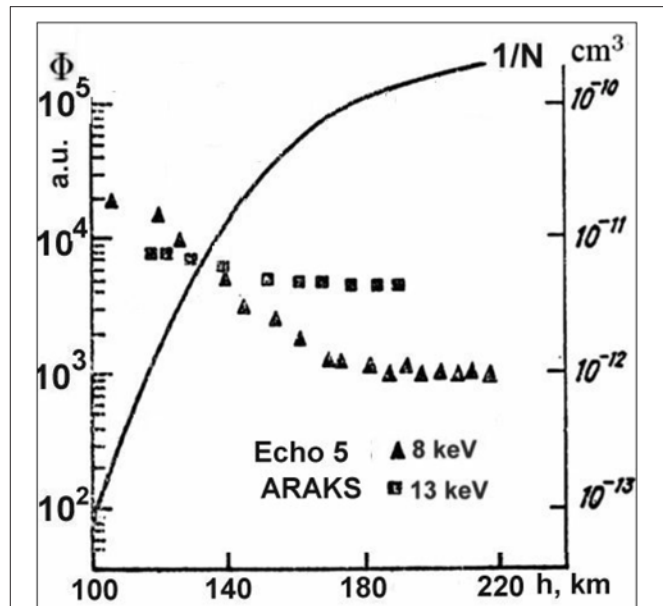


FIGURE 8 | Variation of the scattered beam electrons in the ARAKS (squares) and Echo 5 (triangles) experiments. The solid line shows the model value of N^{-1} in cm^{-3} . After Grandal (1982a) and Arnoldy et al. (1985). Reprinted by permission from COSPAR and American Geophysical Union.

the SPA predicts the decrease by a factor of 30. By the same token, the Polar 5 data reveal (Maehlum et al., 1980c) that an initially-collimated beam significantly broadens over pitch angles in less than one gyroperiod, i.e., $\sim 10^{-6}$ s. Furthermore, in addition to the beam core, a noticeable part of beam electrons is scattered over large pitch angles up to $\sim 180^\circ$. At altitudes 150 to 180 km, the flux of these “halo” electrons varies with the distance, $d_{\perp}$ (in meters), between the mother and daughter payloads as

$$\Phi_{obs} = 3 \cdot 10^7 \exp(-0.07 \cdot d_{\perp}), \quad (5)$$

which is about 10^2 times greater than collisional scattering can produce.

In summary, the observed near-rocket glow, suprathermal electrons, VHF radioemission, and prompt electron echoes, as well as the fine altitudinal structure of AA rays, point to much stronger interaction of injected electrons with the upper atmosphere than provided by electron collisions. Next, a brief survey is presented of the theory of collisionless beam-plasma interaction (BPI) resulting in the beam-plasma discharge (BPD) near beam-emitting payloads and the double-peak structure of artificial and natural aurora rays.

4. BEAM- IONOSPHERIC PLASMA INTERACTION

The upper atmosphere is, in fact, a weakly-ionized plasma, the ionosphere, with the plasma density $n_e \ll N$. Conventionally (e.g., Ichimaru, 1973), ionized gases can be regarded as plasma if their behavior is dominated by the collective response of charged

particles due to the long-range Coulomb force. Charged particles in motion generate electromagnetic fields that affect motion of other particles, thereby making a fast remote response to local perturbations. The distance over which charges in plasma are shielded is the electron Debye radius, $r_D = v_{Te}/\omega_{pe}$, where $\omega_{pe} = 2\pi f_{pe}$ (the Langmuir frequency) and v_{Te} is the electron thermal velocity. The number of particles in the Debye sphere in plasmas is large, $N_D = 4\pi n_e r_D^3 \gg 1$. In collisional plasmas, the plasma frequency is supposed to significantly exceed the collision frequency of plasma electrons, ν_e .

Henceforth, the density, mass, electric (magnetic, B_0) fields, and frequency, $f = \omega/2\pi$, are taken in cm^{-3} , grams, V/m (Gauss), and Hz, respectively. Temperature, $T_{e,i}$, is taken in units of energy, electronvolts (1 eV = 11,605 K). We consider only one singly-charged ion species with the ion-to-electron mass ratio $m_i/m_e = \mu^{-1} = 3 \cdot 10^4$ and $6 \cdot 10^4$ in the F- and E-region ionosphere at altitudes around ~ 200 km and ~ 100 –130 km, respectively.

4.1. Plasma Waves

A symbiotic relationship between plasma particles and fields results in a wide variety of collective motions, i.e., plasma waves. Of those, the most important for the BPI are high-frequency (HF) plasma modes. In a weakly-magnetized ($\omega_{ce} \ll \omega_{pe}$) plasma, their frequencies away from the electron gyroharmonics, $s\omega_{ce}$, are

$$\omega_{\mathbf{k}} = \omega_{pe} \left(1 + \frac{3}{2} k^2 r_D^2 + \frac{\omega_{ce}^2}{\omega_{pe}^2} \sin^2 \alpha \right) \quad (6)$$

with the wavenumber $\mathbf{k}$, the propagation angle $\alpha = \arccos(\mathbf{k} \cdot \mathbf{B}_0/kB_0)$, and $k^2 r_D^2 \ll 1$. It reduces to the Langmuir (L) branch, $\omega_l \approx \omega_{pe} (1 + \frac{3}{2} k^2 r_D^2)$ at $\alpha = 0$ ($k = k_{\parallel}$) and to the upper hybrid (UH) branch, $\omega_{uh} \approx \sqrt{\omega_{pe}^2 + \omega_{ce}^2}$ at $\alpha = \pi/2$ ($k = k_{\perp}$). The spectral energy density of a broad Langmuir spectrum, $\delta E(\mathbf{r}, t) = \sum_{\mathbf{k}} \mathbf{E}_{\mathbf{k}} e^{i\mathbf{k}\mathbf{r} - i\omega_{\mathbf{k}} t}$ is $W_l \approx \sum |\mathbf{E}_{\mathbf{k}}|^2 / 4\pi$.

The other important branch of electrostatic ($\mathbf{E} = -\nabla\phi$) HF waves is the electron Bernstein (EB) mode with frequencies approaching $s\omega_{ce}$ at both $k_{\perp}^{-1} \gg r_{ce}$ and $\ll r_{ce}$ (the electron Larmor radius). Near $k_{\perp} r_{ce} \sim 1$, they have a wide maximum around half-integral gyroharmonics, $\omega_s \approx (s + 1/2)\omega_{ce}$, where the group velocity tends to zero (e.g., Mikhailovskii, 1974). It is worth noting that in the magnetospheric community such waves, routinely observed in the plasma sheet region associated with diffuse aurora (e.g., Ashour-Abdalla and Kennel, 1978; Khazanov, 2011), are called electron cyclotron harmonics (ECH).

It is instructive to recall the well-known wave-particle quantum-mechanical analogy (e.g., Ichimaru, 1973), which helps to understand the wave properties in an inhomogeneous plasma. Since $\omega_{\mathbf{k}}(\mathbf{r}, \mathbf{k})$ is constant in stationary media, we have $d\omega_{\mathbf{k}} = d\mathbf{r} \frac{\partial \omega_{\mathbf{k}}}{\partial \mathbf{r}} \Big|_{\mathbf{k}} + d\mathbf{k} \frac{\partial \omega_{\mathbf{k}}}{\partial \mathbf{k}} \Big|_{\mathbf{r}} = 0$ or

$$\frac{d}{dt} \mathbf{r} = \mathbf{v}_g = \frac{\partial}{\partial \mathbf{k}} \omega_{\mathbf{k}}; \quad \frac{d}{dt} \mathbf{k} = -\frac{\partial}{\partial \mathbf{r}} \omega_{\mathbf{k}} \quad (7)$$

That is, the wave vector changes in such a way that the frequency is preserved. Equations (7) are the Hamilton equations for unit-mass particles (*plasmons*) with the “energy” $\omega_{\mathbf{k}}$ and “momentum”

$\mathbf{k}$, moving with the (group) velocity $\mathbf{v}_g$. It is also convenient to introduce the number of plasmons as $N_{\mathbf{k}} = W_{\mathbf{k}}/\omega_{\mathbf{k}}$.

Let us consider Langmuir plasmons moving in a one-dimensional density depletion (cavity), $n(x) = n_0(1 - |\delta n(x)|)$ centered at $x = x_c$, with the width $L_n \gg k_c^{-1}$ ($k_c = k(x_c)$). Expanding $\omega_{pe}(n)$ gives $\omega_l/\omega_{p0} - 1 \approx \frac{1}{2} \delta n(x) + \frac{3}{2} k^2(x) r_D^2$, so the Hamiltonian can be represented as the sum of the kinetic energy, $\varepsilon_{\mathbf{k}} = \mathbf{k}^2/2$, and potential, $U(x) = 2\delta n(x)/3r_D^2$. As it follows from (7), $\frac{d}{dt} \mathbf{k} = \mathbf{v}_g \frac{\partial}{\partial \mathbf{r}} \mathbf{k} > 0$ (“acceleration”) and $\frac{d}{dt} \mathbf{k} < 0$ (“deceleration”) for plasmons moving toward and from the cavity’s center, respectively. Evidently, the cavity plays a role of the potential hole for plasmons that are trapped inside if the Hamiltonian is negative, i.e., $-U(x_c) > \varepsilon_{\mathbf{k}}(x_c)$ or

$$|\delta n(x_c)| > 3k_c^2 r_D^2, \quad (8)$$

and move freely otherwise.

4.2. Resonance Wave-Particle Interaction

Let us consider plane Langmuir waves, $\mathbf{E}(\mathbf{r}, t) = \mathbf{E}_{\mathbf{k}} e^{i\mathbf{k}\mathbf{r} - i\omega_{\mathbf{k}} t}$, moving along the magnetic field ($\parallel \mathbf{z}$) much faster than the bulk of electrons, i.e., $v_{Te} \ll \omega_{\mathbf{k}}/k_z = v_{ph}$ (the phase velocity). Thus, only a small group of fast particles can be in resonance with the waves, i.e., $v_{\parallel} = v_{res} = (\omega_{\mathbf{k}} - s\omega_{ce})/k_{\parallel}$, with $s = 0, \pm 1, \pm 2$, etc. It is called the Cherenkov resonance at $s = 0$ and the cyclotron resonance otherwise. It is instructive to give an example of the resonance wave-particle interaction in isotropic plasmas. The equation of motion of an electron with the unperturbed velocity $\mathbf{v}_0 = \text{const}$ reads

$$\frac{d}{dt} \mathbf{r} = \mathbf{v}_0 + \mathbf{v}_E; \quad \frac{d}{dt} \mathbf{v}_E = \frac{-e}{m_e} \mathbf{E}_{\mathbf{k}} \exp(i\mathbf{k}\mathbf{r} - i\omega_{\mathbf{k}} t) \quad (9)$$

(e is the elementary charge). Linearizing in the wave field yields the quiver velocity

$$\mathbf{v}_E \approx -\frac{ie}{m\omega'_{\mathbf{k}}} \mathbf{E}_{\mathbf{k}} e^{-i\omega'_{\mathbf{k}} t} \quad (10)$$

where $\omega'_{\mathbf{k}} = \omega_{\mathbf{k}} - \mathbf{k}\mathbf{v}_0$ is the Doppler-shifted wave frequency. The work of the electric field, $\propto \text{Re}(\delta \mathbf{v}) \cdot \text{Re}(\mathbf{E})$, upon non-resonance particles, $\omega'_{\mathbf{k}} \neq 0$, vanishes as the positive and negative contributions on average cancel each other. For resonance particles, $\omega'_{\mathbf{k}} \rightarrow 0$, the linearization procedure is invalid, and Equation (9) must be solved explicitly. Taking $\mathbf{k}$ along the z axis, in the reference system of the wave, $z' = z - v_{ph} t$, yields the equation of motion in a periodic wave potential

$$\frac{d^2}{dt^2} \xi + \omega_{tr}^2 \sin \xi = 0 \quad (11)$$

Here $\xi = kz'$ is the relative phase, $\phi_{\mathbf{k}, \omega}$ is the amplitude of the potential, and $\omega_{tr} = \sqrt{\frac{e}{m_e} k^2 |\phi_{\mathbf{k}}|}$. Nearly at the resonance, $\xi \ll 1$, Equation (11) reduces to a classic (mechanical) oscillator with $\xi \approx \xi_0 \cos(\omega_{tr} t)$. That is, particles in the resonance zone, $v_z^{(-)} < v_z < v_z^{(+)}$, where

$$v_z^{(\pm)} = v_{ph} \pm v_{tr} \text{ and } v_{tr} = \sqrt{\frac{e}{m_e} |\phi_{\mathbf{k}}|} \ll v_{Te}, \quad (12)$$

are trapped and oscillate around the minimum potential energy, maintaining the constant phase relative to the wave for many wave periods (e.g., Galeev and Sagdeev, 1979). The particles exchange energy with the wave

$$\Delta\varepsilon = 2m_e v_{ph}(v_{ph} - v_z) \quad (13)$$

after reflection from a moving potential barrier. That is, slow ($v_z^{(-)}$) particles are “kicked” along by the wave and gain energy, while fast ($v_z^{(+)}$) particles push on the wave and lose energy. The net energy exchange is defined by the difference, $F_0(v_z^{(+)}) - F_0(v_z^{(-)}) \approx 2v_{tr} \left. \frac{\partial F_0}{\partial v_z} \right|_{v_{ph}}$, between the population

of the two groups with the distribution function $F_0(\mathbf{v})$. It is negative in a Maxwellian plasma, so the waves are damped. This resonance, collisionless damping is named the Landau damping. In a magnetoactive plasma, oblique waves are subjected to the cyclotron damping as well. Therefore, in thermal equilibrium waves exist at the thermal noise level, $W \sim n_e T_e / N_D$.

In non-equilibrium plasma, the population of fast particles can dominate in a certain velocity domain so that waves in resonance with these particles will gain energy and grow. This process (inverse Landau or cyclotron damping) is called plasma instability.

4.3. Beam (Bump-in-Tail) Instability

Let a collimated, $u = v_{\parallel} \gg v_{\perp}$, “warm” and tenuous, $1 > \Delta u_b / u_b \gg (n_b / n_e)^{1/3}$, electron beam with the beam density, $n_b \ll n_e$, and the mean energy, $\varepsilon_b \gg T_e$, precipitates into the ionosphere along the magnetic field. Such distribution function is known as a “bump-in-tail,” meaning the tail of the whole electron distribution with the bulk, secondary, and beam electrons. As $\partial F_b / \partial u > 0$ at $u_b - \Delta u_b < u < u_b$, Langmuir waves with phase velocities within this range will grow at a rate (e.g., Mikhailovskii, 1974)

$$\Gamma_r = \frac{dW_r}{W_r dt} = \gamma_b - \nu_e \approx \omega_{pe} \frac{\pi n_b}{n_e} \left(\frac{u_b}{\Delta u_b} \right)^2 - \nu_e - \gamma_{nl} \quad (14)$$

Here W_r is the spectral energy density of the beam-resonant waves, $k_{\parallel} = k_r$, γ_{nl} is the rate of spectral transfer due to nonlinear mode coupling (see below), and $\nu_e \approx 10^{-7} T_e^{5/6} N$ at $T_e < 0.4$ eV and $\approx 10^{-7} T_e^{1/2} N$ at $0.4 < T_e \leq 100$ eV below ~ 200 km. Calling for $\Gamma_r > 0$ at $\gamma_{nl} = 0$ gives the limiting neutral density

$$N < N_{\max} \approx 1.5 \cdot 10^{12} \frac{\gamma_b}{\omega_{pe}} T_e^{-5/6} n_e^{1/2} \text{ cm}^{-3} \quad (15)$$

Taking $n_e \sim 10^5 \text{ cm}^{-3}$, $T_e \sim 0.1$ eV, $n_b \sim 1 \text{ cm}^{-3}$, and $\Delta v_b / v_b \sim 0.1$ yields for the standard neutral atmosphere that the instability develops at altitudes above $h_{\min} = h(N_{\max}) \sim 105$ km.

The wave excitation goes at the expense of the beam energy leading to widening of the beam distribution and thus decrease of γ_b . In a steady state, the dissipation rate is determined by the energy flux balance

$$n_b \frac{\partial}{\partial z} \langle u_b \varepsilon \rangle_b = - \frac{\partial}{\partial z} v_g W_r \approx - \Gamma_r W_r \quad (16)$$

where $\langle \dots \rangle_b$ means averaging over the beam distribution. The beam speed greatly exceeds the wave group velocity, $v_g \approx 3T_e / m_e u_b$, so in the absence of nonlinear interactions W_r is greater than $n_b \varepsilon_b$ by a factor of $v_b / v_g \sim \varepsilon_b / T_e$.

The relaxation length, l_{rel} , is a distance from the beam entry into the plasma over which the instability stabilizes at some $\Delta u_b(l_{rel}) = \Delta u_{\infty}$. It can be estimated from Equation (16) as follows

$$l_{rel} \sim \Lambda \frac{\Delta u_{\infty}}{\gamma_b^{\infty}} \frac{n_b \varepsilon_b}{W_r} \quad (17)$$

Here $\Lambda \sim \ln(N_D) \sim 10$ is a numerical coefficient accounting for the growth of the waves from the thermal level. In order to determine the level of the beam-excited waves, W_r , one should allow for nonlinear wave interactions resulting in the energy transfer from the resonance region at a rate γ_{nl} . The efficiency of this nonlinear damping is governed by the parameter of nonlinearity, $w = W / n_0 T_e \ll 1$.

4.4. Nonlinear Effects

Disregarding wave-wave interactions and using the statistical description of waves with random phases, results in the quasilinear approximation (Vedenov et al., 1962). It is valid only for very tenuous beams irrelevant to our problem. The next step in w is taking account of the induced scattering, in which electrons interact with beats of different randomly-phased modes (e.g., Galeev and Sagdeev, 1979). This weak turbulence (WT) approximation is valid until the beam density exceeds

$$n_b^{(th)} \approx 0.1 \mu \frac{T_i}{T_e} n_e \sim (1 - 3) \cdot 10^{-6} n_e \quad (18)$$

(Galeev, 1975; Papadopoulos, 1975; Galeev et al., 1977). Then, the beam relaxation is described in terms of strong Langmuir turbulence (SLT).

The SLT regime is inherently tied to the tendency of Langmuir plasmons to accumulate inside density depletions or cavities, $\delta n_s = n_e / n_0 - 1 < 0$. Namely, the waves of wavelengths $\sim k^{-1}$ are trapped if $|\delta n_s| > 3k^2 r_D^2$ (Equation 8). The resulting excess of the wave pressure, $\delta W_l \approx W_l |\delta n_s| / 3k^2 r_D^2$, exceeds the thermal pressure imbalance, $\delta p_e = n_0 T_e \delta n_s$, if $w_l > w_{th} = 3k^2 r_D^2$. In this case, the ponderomotive force, $-\nabla \delta W_l$, pushes plasma out of the cavity, which further deepens and traps yet more plasmons in a positive-feedback loop. As a result, initial modulations grow with time. At $W_l / n_0 T_e \gg W_{th}$, the growth rate of this modulational instability (MI) is

$$\gamma_{mi}(W_l) \approx \omega_{pe} \sqrt{\frac{\mu w_l}{3}} \quad (19)$$

As trapping inside a cavity leads to strong correlation of the wave phases, such that the WT condition of random phases (e.g., Galeev and Sagdeev, 1979) is violated, this regime has been termed strong Langmuir turbulence (SLT).

Cavities with trapped strongly-correlated Langmuir oscillations are termed cavitons and subjected to collapse (Zakharov, 1972). Their evolution depends on the dimension,

d , and can be understood from simple arguments (e.g., Sagdeev, 1979). The conservation of the plasmons' number in a cavity of the size, l , yields $|\mathbf{E}^2(t)| \propto l^{-d}(t)$. The wavelengths of the trapped plasmons are also of the order of l , i.e., $k \sim 1/l$. The trapping condition yields $l^{-1} \sim k \propto |\delta n_s|^{1/2}$, indicating that a deepening cavity narrows, *collapses*, as time progresses. Since $|\delta p_e| \propto T_e |\delta n_s| \propto l^{-2}$ and $|\mathbf{E}^2| \propto l^{-d}$, the thermal pressure will ultimately balance the HF pressure for $d = 1$, thus forming one-dimensional cavitons. In two (three) dimensions, the speed of collapse persists (accelerates) with time.

In a weakly-magnetized plasma, $\omega_{ce} \ll \omega_{pe}$, the cavitons at $|\delta n_s| \ll \omega_{ce}^2/\omega_{pe}^2$ are pancake-like, with the dimensions

$$l_{\parallel} \sim k_{\parallel}^{-1} \sim r_D |\delta n_s|^{-1/2} \sim l_{\perp} \frac{\omega_{pe}}{\omega_{ce}} |\delta n_s|^{1/2} \quad (20)$$

(e.g., Rowland et al., 1981b; Shapiro and Shevchenko, 1984; Robinson, 1997). The basic signatures of the SLT development in various beam-plasma systems have been observed in laboratory experiments (e.g., Cheung et al., 1982; Karfidov and Lukina, 1997; Robinson, 1997; Vyacheslavov et al., 2002).

The phase velocity of plasmons in collapsing cavitons, $\sim \omega_{pe}/k(t)$, decreases with time, so eventually plasmons are absorbed by plasma electrons due to Landau and transit-time damping. As a result, a small group of suprathermal electrons gains energy, while the HF pressure in the caviton drops and collapse is arrested due to the wave energy “burnout.” Ultimately, a dynamic equilibrium is reached between the pumping energy into cavitons in the long-scale *source* region, $k_L \leq r_D^{-1} \sqrt{W_L/3}$, and short-scale transfer by collapsing cavitons (W_{cav}) into the absorption interval, $k \geq k_a$ (e.g., Galeev et al., 1977). The energy density in the source region, $W_L > W_{cav} > W_a = \int_{k \geq k_a} W_k dk$, comprises the MI-excited (non-trapped) long-scale waves.

In collisionless isothermal plasmas one gets at $\gamma_b \gg \mu \omega_{pe}$ (Galeev et al., 1977)

$$w_r \approx 3 \left(\frac{\mu \gamma_b}{\omega_{pe}} \right)^{1/2} \ll w_L \approx 3 \frac{\gamma_b}{\omega_{pe}} \quad (21)$$

In collisional plasmas, Equations (18) and (21) hold for $v_e/\omega_{pe} < 3T_e/2\varepsilon_b$. In the opposite case, the MI threshold is $w_{th}^{(c)} \approx 2v_e/\omega_{pe}$ and the threshold beam density becomes (Volokitin and Mishin, 1979)

$$n_b^{(c)} \approx 0.1 n_e \mu \frac{T_i}{T_e} \frac{v_e}{\omega_{pe}} \frac{\varepsilon_b}{T_e} \quad (22)$$

If $W_L \gg W_r > W_{th}^{(c)}$, the MI growth rate is of the same order as γ_{mi} (19). However, as follows from eq. (21), $\gamma_{mi}(W_L) < v_e$ at

$$v_e > v_* \approx (\mu \gamma_b \omega_{pe})^{1/2} \quad (23)$$

collisional damping is faster than collapse. Since the nonlinear transfer rate due to collapse reduces, the level of plasma waves in this “collisional SLT” regime increases over that

in Equation (21) and becomes (Volokitin and Mishin, 1979; Mishin and Telegin, 1989)

$$w_r^{(*)} \approx \frac{3v_e}{\omega_{pe}} \ll w_L^{(*)} \approx \frac{3}{\mu} \left(\frac{v_e}{\omega_{pe}} \right)^2 \quad (24)$$

Now the relaxation length can be estimated using Equations (17), (21), and (24)

$$l_b \sim v_b \frac{\varepsilon_b}{T_e} \left(\frac{\Delta u_{\infty}}{v_b} \right)^3 \begin{cases} v_*^{-1} \text{ at } v_e < v_* \\ v_e^{-1} \text{ at } v_e > v_* \end{cases} \quad (25)$$

In addition, in the presence of short-scale, $q \gg k_L$, density irregularities, $\delta n_s(\mathbf{r}) = \sum_q \text{Re}(\delta n_q \exp(iqz))$, long-scale plasmons are transferred into the short scales via conversion, $L_{k_L} + \delta n_s \rightarrow L'_q$ (Galeev et al., 1977). For a wideband random-phase oscillations, $|\delta n_q| < 3q^2 r_D^2$, the conversion rate is

$$\gamma_{conv} \approx \sum_q v_l(\mathbf{q}) \frac{|\delta n|^2_{\mathbf{q}}}{36q^4 r_D^4} \ll v_l(\mathbf{q}) < 3\omega_{pe} q^2 r_D^2 \quad (26)$$

Here $v_l(\mathbf{k})$ is the total (collisional + Landau) damping rate and $|\delta n|^2_{\mathbf{q}}$ is the phase-averaged spectral energy. The conversion process dominates at $\gamma_{conv} > \gamma_{mi}(w_L)$. Applying a similar procedure to the electromagnetic version of the Zakharov equation (e.g., Shapiro and Shevchenko, 1984), one can describe resonant scatter of radio waves on ion density oscillations (e.g., Mishin et al., 1992).

The value of T_e in Equation (25) is the average energy of the bulk electrons, $T_e = T_h$, heated by the beam-excited turbulence at a rate

$$\tau_{heat}^{-1} \approx \frac{2}{3} v_e(T_e) W_l / n_e \quad (27)$$

At altitudes below ~ 130 km the growth of T_e is limited mainly by inelastic losses, $v_{il}(T_e) = \delta_{il}(T_e) v_e(T_e)$. The coefficient of inelastic losses at $h = 150$ – 180 km, calculated using the Majeed and Strickland (1997) tabulations, is $\delta_{il}(T_e) \approx (0.1 \rightarrow 1.5) \cdot 10^{-2} \rightarrow 0.1 \rightarrow 0.3 \rightarrow 0.1$ at $T_e \approx (0.2 \rightarrow 0.45) \rightarrow 0.6 \rightarrow 5 \rightarrow 10$ eV. Assuming the ambient temperature $T_0 = 0.2$ eV and $W_l/n_e T_0 \sim 10^{-3}$ – 10^{-2} in a steady state, one gets the temperature of heated electrons $T_h \approx T_0 (1 + 2W_l/3\delta_e(T_h)) \approx 0.3$ – 0.5 eV.

4.5. Acceleration of Suprathermal Electrons

The SLT acceleration of suprathermal (tail) electrons is, probably, the most important consequence of the BPI for artificial and natural aurora. In brief, the short-scale, $k > k_a$, plasmons in collapsing cavities are absorbed by a small group of plasma electrons. Their distribution function along the magnetic field can be found from the kinetic equation

$$\frac{d}{dt} F_t = \frac{\omega_{pe}^2}{m_e n_0} \frac{\partial}{\partial u} \frac{W_{\omega_k/u}}{u} \frac{\partial F_t}{\partial u} \equiv \frac{\partial}{\partial u} D(u) \frac{\partial F_t}{\partial u} \quad (28)$$

As a result, a suprathermal tail, $\varepsilon \gg T_e$, is formed, with a power-law distribution

$$F_t(\varepsilon) = \frac{n_t}{8\pi v_{\min}^3} \left(\frac{\varepsilon}{\varepsilon_{\min}} \right)^{-p_t} \quad \text{at } \varepsilon_{\max} \geq \varepsilon \geq \varepsilon_{\min} \quad (29)$$

In a Maxwellian isotropic plasma ($F_0 = F_M$), theoretical estimates give the spectral index $p_t \approx 7/4 - 9/4$ and the minimum energy, $\varepsilon_{\min}^{(M)} \sim 10T_e$; while the matching condition, $F_t(\varepsilon_{\min}) = F_0(\varepsilon_{\min})$, yields $n_t^{(M)} \sim 10^{-4}n_e$ (Galeev et al., 1977; Pelleiter, 1982). One-dimensional numerical simulations yield $p_t \approx 3/2$ (Galeev et al., 1983; Wang et al., 1997). **Figure 9** illustrates the wave spectrum and electron distribution in the developed strong Langmuir turbulence.

Substituting the accelerated flux $\Phi_t(\varepsilon)$ into Equation (2) with $\sigma_\lambda = \sigma_{ion}$ gives the ionization rate, which can be approximated as (Mishin and Telegin, 1989)

$$q_{ion}^{(t)} \approx n_b v_e(T_e) \frac{3T_e}{\varepsilon_{ion}} \left(\frac{u_b}{\Delta u_b} \right)^4 \quad (30)$$

A remark is in order. In the presence of the ambient (“seed”) suprathermal population, such as secondary electrons with the distribution $F_s(\varepsilon) \gg F_M(\varepsilon)$ at $\varepsilon \geq \varepsilon_{\min}^{(M)}$, the absorption rate increases and collapse is arrested at greater scales than in a Maxwellian plasma. That is, $\varepsilon_{\min}^{(s)} \gg \varepsilon_{\min}^{(M)}$ and thus many more energetic electrons are accelerated (Mishin and Telegin, 1986). Another effect of electron-neutral collisions is that $\varepsilon_{\max}$ depends on inelastic losses of the accelerated electrons, $\sim v_{ion}(\varepsilon)$, that maximize at $\varepsilon_m \sim 100$ eV. Evidently, as soon as the rate of inelastic losses, $v_{il}(\varepsilon_L)$ at $\varepsilon_L \approx 10^4 W_L/n_e < \varepsilon_m$, exceeds the acceleration rate, $m_e D(u_L)/8\pi \varepsilon_L$, the expansion of the tail, $F_t(\varepsilon) \propto \varepsilon^{-p_t}$, stops (cf. **Figure 6B**). In the SLT regime (24), it

occurs at $N > N_m \approx 10^8 \sqrt{n_e} \text{ cm}^{-3}$ (Volokitin and Mishin, 1979; Mishin, 2010).

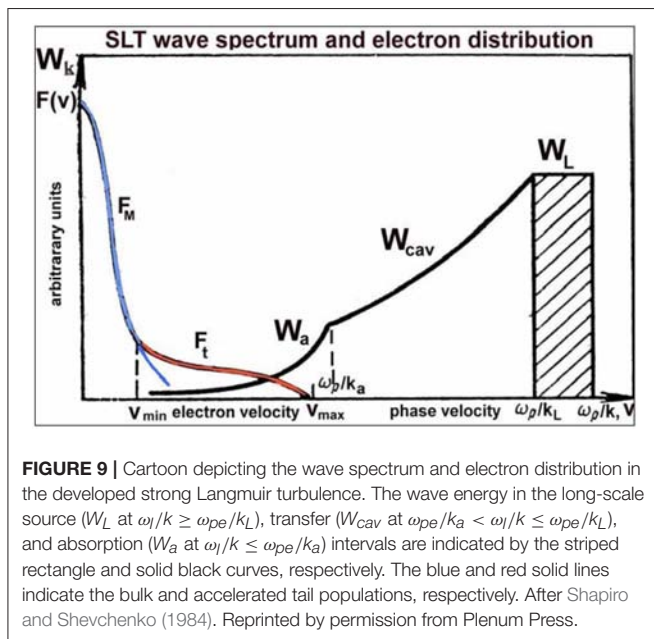
4.6. SLT Acceleration in High-Power Radio Wave Experiments

It seems relevant to briefly discuss the SLT acceleration of suprathermal electrons producing artificial aurora and ionization in active space experiments with high-power radio waves (Mishin and Pedersen, 2011; Eliasson et al., 2012, 2015; Mishin et al., 2016). Here Langmuir waves are parametrically driven by ordinary (O) pump waves near the reflection altitude, h_0 , where the pump frequency, f_0 , matches the local plasma frequency, $f_{pe}(h_0)$ (the interested reader is referred to Streltsov et al., 2018 review). **Figure 10** exemplifies the results of the Eliasson et al. (2012, 2015) full-wave one-dimensional simulations with various input pump amplitudes, E_{in} , for radio beam pointings at and between the geographic zenith (vertical, V) and the geomagnetic field direction (magnetic zenith, MZ) at the High-frequency Active Auroral Research Program (HAARP) facility at Gakona, Alaska, USA. The simulation details are given in Eliasson et al. (2012, 2015). Note only that the nonlinearity parameter, $w \propto E_{in}^2/8\pi n_e T_e$, exceeds the threshold for the SLT regime to develop.

Figures 10a,b illustrates the development of the longitudinal (Langmuir) electric field, E_z , and cavitons, δn_s , for vertical and MZ injections with $E_{in} = 1$ and 2 V/m. Omitting the features of their spatial distribution related to the radio wave propagation, we point out that in 1–2 ms the initial (Airy) structure starts breaking into small-scale turbulence. In saturation, solitary wave packets in the SLT region are trapped in density cavitons. It is seen that the SLT region of the altitudinal extent, l_{LT} , is sandwiched between the WT regions with turbulent electric fields but without cavitons. The appearance of the electromagnetic waves near h_0 at MZ is due to $O + \delta n_k \rightarrow O'$ conversion on short-scale ion density oscillations (Equation 26) in the SLT region. **Figures 10c,d** show the turbulent electric fields, δE_{sat} , and the electron energy distribution, $F_t(\varepsilon) = F_t(u)du/d\varepsilon$, at the end of the simulation runs.

At each injection angle, the simulated spectral width, Δk , as well as δE_{sat} and l_{LT} , increase with E_{in} . These factors and the input value of T_e lead to considerable differences in $F_t(\varepsilon)$. Overall, the main part of $F_t(\varepsilon)$ at $\varepsilon_{\max} \geq \varepsilon \geq \varepsilon_{\min}$ can be fitted by a power law, $F_t(\varepsilon) \propto \varepsilon^{-p_t}$, with the relative tail density, $10^4 n_t/n_e$, between 2 and 6 and p_t between ≈ 1.5 and 2. The maximum energy, $\varepsilon_{\max}$, at altitudes $h \geq 170$ km depends mainly on the transit time, $\tau_{\parallel} \sim l_{LT}/u_{\max}$, while inelastic losses dominate below 170 km. In particular, $F_t(\varepsilon)$ at 10.5° S is more enhanced than at vertical due to greater l_{LT} . At MZ for $E_{in} = 2$ V/m, $F_t(\varepsilon)$ is close to that at vertical for $E_{in} = 1.5$ V/m (Eliasson et al., 2012, **Figure 8**). This explains the differences in the patches (layers) of artificial plasma descending from the initial interaction altitude at various input parameters shown in **Figure 10e**.

The downward propagation of the artificial plasma produced by the accelerated electron tail is due to the fact that the electrons propagating along the geomagnetic field create the new plasma resonance condition for the incident radio wave below the initial resonance. This way, an ionizing wavefront created by



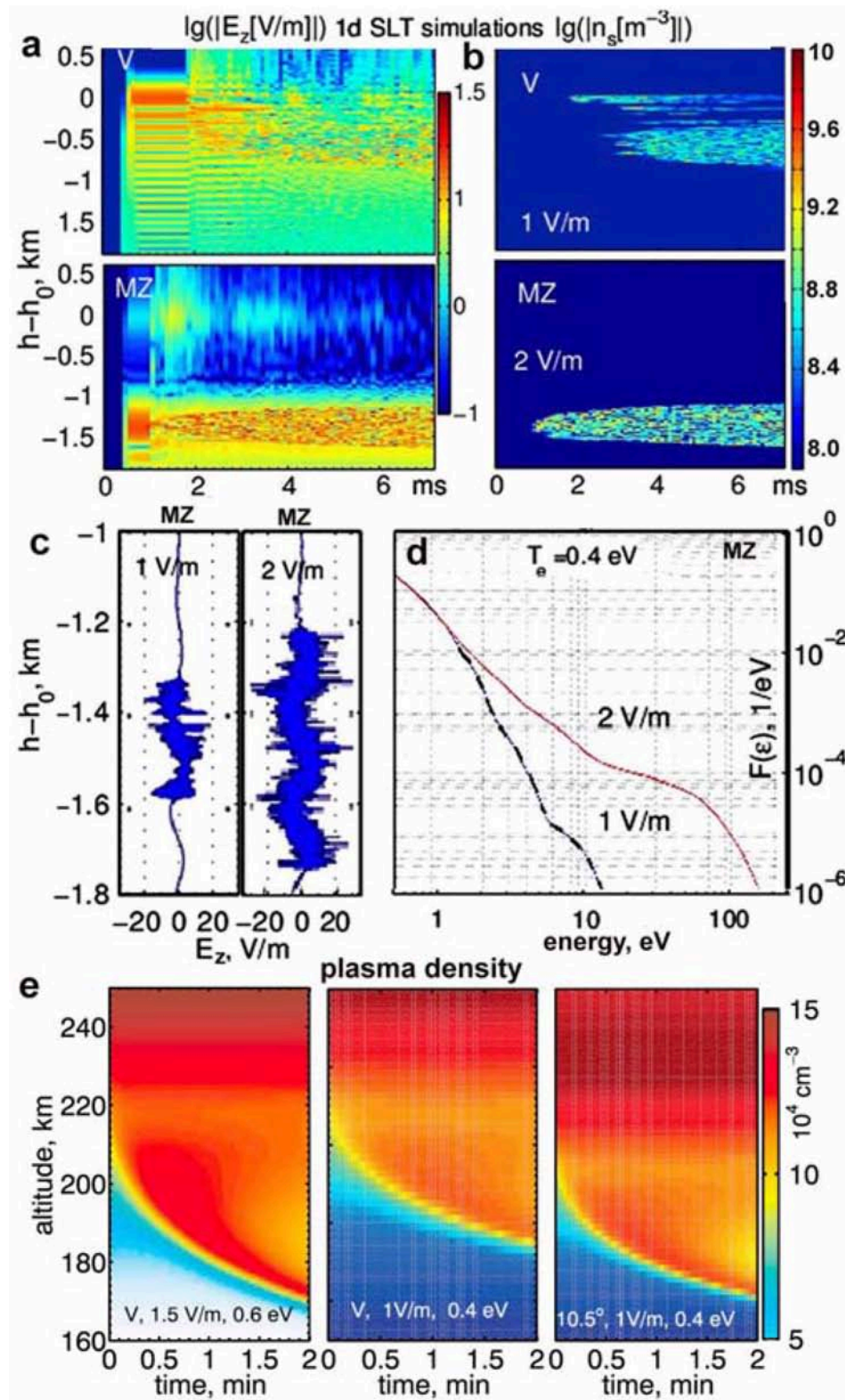


FIGURE 10 | Time-vs.-altitude color-coded plots of (a) $\log_{10}(|E_z[\text{V/m}]|)$ and (b) $\log_{10}(|\delta n_s[\text{m}^{-3}]|)$ for $E_{in} = 1$ and 2 V/m at V and MZ, as indicated in plots; (c) turbulent electric fields, E_z , (d) the distribution function, $F_t(\epsilon) = F_t(u)du/d\epsilon$; and (e) time-vs.-altitude plots of the plasma density at V and 10.5° S calculated for various E_{in} and T_e indicated in frames. Adapted from Eliasson et al. (2012, 2015). Reprinted by permission from the American Geophysical Union.

the SLT-accelerated electrons is formed (Mishin and Pedersen, 2011; Eliasson et al., 2012). Furthermore, the presence of the ambient suprathermal population (photoelectrons) facilitates the SLT acceleration (Mishin et al., 2004, 2016; Eliasson et al., 2018), quite similar to the effect of secondary electrons in auroral plasmas (Mishin and Telegin, 1986). The consequences of the photoelectrons in the sunlit ionosphere are the decreased threshold, the greater downward speeds, and the decay of the persistent artificial ionization at the terminal altitude after sunlit-to-dark transition.

Next we explore the effect of transition from the “collisionless” to “collisional” SLT regime at $v_e > v_*$ (23).

4.7. Plasma-Turbulence Layer

At $v > v_*$ (23), the plasma turbulence level increases as $W_L^{(*)} \sim v_e^2$ (24). Thus, the condition $v_e(h_*) = v_*$ defines the upper boundary, h_* , of the layer of enhanced plasma turbulence termed the Plasma Turbulence Layer or the PTL (Mishin and Telegin, 1989; Mishin et al., 1989). Balancing the heating rate (27) by inelastic losses gives

$$\delta_{il}(T_e) \approx 3 \frac{\gamma_b}{\omega_{pe}} \begin{cases} 1 & \text{at } v_e < v_* \\ (v_e/v_*)^2 & \text{at } v_e > v_* \end{cases} \quad (31)$$

As follows from (31), for $\gamma_b/\omega_{pe} \sim (0.3-1)10^{-3}$ the temperature T_e reaches $\approx 0.3-0.45$ eV in the PTL and does not exceed ≈ 0.25 eV above h_* . Since $q_{ion}^{(t)}$ (30) increases with T_e , the auroral luminosity is peaked inside the PTL and greatly enhanced over the collisional (SPA) limit, $q_b \sim n_b v_b$. This regime continues until $v_e(h)$ increases to $v_e(h_*) = v_* \approx 0.3\mu^{1/2}\omega_{pe}$. Below h_* , deep cavities are not created and $w_{r*}^* \sim w_{L*}^* \sim \mu^{-1} \left(\frac{\gamma_b}{\omega_{pe}} \right)^2$, so collapse and concomitant suprathermal electron acceleration are inhibited. For typical auroral beam-plasma parameters, $n_b/n_e \sim 10^{-5}$ and $\Delta u/u_b \sim 0.2$, the value of v_* is of the order of $3v_*$. This means that for the standard neutral atmosphere scale height, $H_N \sim 5$ km, the overall PTL thickness is of the order of 5 km.

Figure 11 illustrates a scenario (Mishin et al., 1981; Mishin and Telegin, 1989; Mishin, 2010) of the formation of a double-peaked ionization/luminosity profile of auroral rays due to the PTL (cf. **Figure 4**). The lower peak ($q_{ion}^{(c)}$) is caused by the collisional ionization of neutral gas by the primary (beam) electrons, while the upper layer, i.e., the PTL, is due to accelerated suprathermal electrons. For a few keV beams, both peaks will overlap, so that the observer would see only one thick layer with a sharp upper boundary.

This scenario agrees well with the altitude profile of artificial auroral rays far from the beam-emitting rocket (**Figure 4**) and natural auroral rays (section 6). Note that artificial beams expand across B_0 due to collisional diffusion, so that the beam density far beneath the rocket in Zarnitza 2 was $\sim 1 \text{ cm}^{-3}$ (Izhovkina, 1978), i.e., close to that of natural beams. Also, one should bear in mind that beams lose only $\leq 20\%$ of their energy in the BPD region, acquiring the velocity scatter $\Delta u/u_b \sim 0.1-0.2$ (Mishin and Ruzhin, 1980a, 1981).

A remark is in order. So far our consideration was limited to collisionless interaction of a warm, tenuous (bump-in-tail) beam

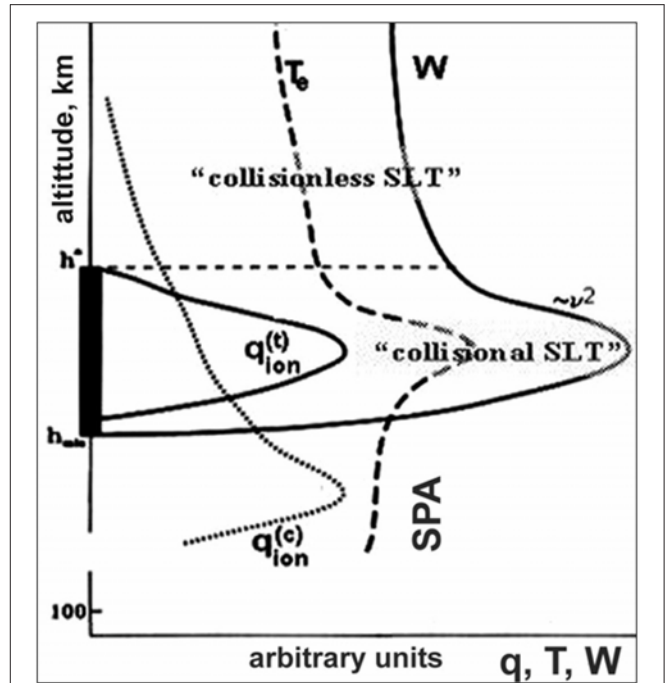


FIGURE 11 | Schematic representation of the Plasma Turbulence Layer (PTL) indicated by the black rectangle and a double-peaked ionization profile. The dotted curve presents the collisional (SPA) ionization profile, $q_{ion}^{(c)}$. The dashed curve shows the electron temperature profile of the heated plasma. The solid curves show the altitudinal profiles of the wave energy density, W , and ionization rate of the accelerated electrons, $q_{ion}^{(t)}$. The dashed horizontal line indicates the PTL upper boundary. After Mishin (2010). Reprinted by permission from the American Institute of Physics.

pertinent to natural and artificial auroral rays. With regards to the near-rocket plasma, this approximation becomes applicable after the plasma density significantly increases over the background in the beam-plasma discharge, which is described below. In addition to the beam instability, highly oblique EB/ECH waves can also be excited for electron injections at large pitch angles due to the non-equilibrium beam distribution over transverse velocities (e.g., Mikhailovskii, 1974). In the following, both the Langmuir and EB/ECH-related processes are outlined.

5. BEAM-PLASMA DISCHARGE

It has long been known that injection of powerful electron beams in neutral gas may result in the avalanche-like ionization accompanied by strong plasma oscillations (e.g., Getty and Smullin, 1963; Vlasenko et al., 1976a; Bernstein et al., 1978, 1979; Szuszczewicz et al., 1979). This phenomenon is termed the Beam-Plasma Discharge (BPD) to emphasize the chief role of the beam-excited waves. As any discharge, BPD develops under certain breakdown conditions that can be readily obtained in a simplified form.

5.1. Qualitative Considerations

At first glance, BPD can be treated similar to the classical high-frequency (HF, $\omega_0 \gg v_e$) discharge (e.g., MacDonald, 1966), in

which the pump waves are excited by the injected beam. In the HF discharge theory, the key parameter is the pump threshold amplitude, E_{hf} , as a function of the applied frequency (ω_0), gas density (N), and size (L) of the discharge gap. The Townsend criterion requires that an electron must acquire enough energy to produce at least one ionizing impact before disappearing from the gap due to recombination and diffusion. Replacing W_l in Equation (27) by $W_{hf} = (\omega_{pe}/\omega_0)^2 |E_{hf}|^2 / 4\pi$ gives the time of the electron heating up to the energy $\varepsilon_h \geq \varepsilon_{ion}$. In HF discharges $\tau_{hf} \propto \varepsilon_h v_e^{-1}(\varepsilon_h) |E_{hf}|^{-2}$ is much greater than τ_{ion} , so the breakdown criterion reduces to

$$\tau_{hf} < \tau_{loss} \quad (32)$$

In the BPD, the pump waves are generated via the beam instability. Therefore, the first step is to find the critical beam density, $n_b^{(cr)}$, necessary for the instability to develop. Let n_{cr} be the minimum plasma density for which the instability can develop at given N and γ_b , similar to the reverse condition (15). As the pre-discharge ionization is created solely by the beam electrons at a rate $q_b \sim v_b n_b$ ($v_b = v_{ion}(\varepsilon_b)$), the instability criterion follows from the ionization balance (Lebedev et al., 1976)

$$n_b > n_b^{(cr)} \sim n_{cr}(\nu_b \tau_{loss})^{-1} \quad (33)$$

Conditionally, the BPD problem may be divided into the “beam” and “discharge” parts. The former considers the beam relaxation to find the input parameters, i.e., the excited wave spectrum and the size of the discharge “gap,” for the latter. Then, the energization (heating and acceleration) of plasma electrons with subsequent ionization is calculated. The solution of the first part critically depends on the plasma density and temperature that vary during the breakdown. Nonetheless, the two parts can be considered independently because the BPI timescale is much shorter than τ_{ion} . That is, the beam relaxation in the course of the breakdown comes about as in a stationary plasma with the density $n_e = n_e(t)$ and temperature $T_h(t)$.

The avalanche starts at some point, t_* , when the ionization by plasma electrons, $q_{ion} = v_{ion}(T_h)n_e(t_*)$, becomes greater than $\nu_b n_b$. Therefore, the BPD threshold is the Townsend condition (32) with the left-hand side replaced by t_* . If suprathermal electrons are disregarded, t_* is of the order of the heating time of the bulk electrons, τ_{heat} . The latter can be estimated from Equation (27) substituting W_l in a generic form $W_l = a_b n_b \varepsilon_b$, where a_b is a coefficient defined by the beam relaxation regime. This yields the criterion for a self-sustained discharge

$$\begin{aligned} \tau_{heat} &\sim 3v_{ion}^{-1}(T_h) \frac{n_{cr}\varepsilon_{ion}}{a_b n_b \varepsilon_b} < \tau_{loss} \\ \text{or } n_b &> n_b^{(d)} \sim (\varepsilon_{ion}/a_b \varepsilon_b) n_b^{(cr)} \end{aligned} \quad (34)$$

Here we accounted for the fact that $\nu_b \sim 3 \cdot 10^{-8} N \text{ s}^{-1}$ for 3–10 keV electrons is of the order of $\frac{1}{3} v_{ion}(T_h)$ at $T_h \geq 1 \text{ eV}$ (e.g., Majeed and Strickland, 1997). The condition (34) is satisfied automatically at $n_b > n_b^{(cr)}$ if $a_b > \varepsilon_{ion}/\varepsilon_b \sim 10^{-3}$, which is valid already for rather tenuous (“weak”) beams, $1 \gg n_b/n_{cr} > 10^{-6}$.

For injections at $h > 100 \text{ km}$, the requirement (33) becomes obsolete as the beam instability develops already in the ambient ionospheric plasma of the density n_0 . A finite cross-section of injected beams reduces the growth rate but not quenches the instability (e.g., Alekhin et al., 1972). Furthermore, the values of τ_{heat} and τ_{loss} at the beginning of the BPD significantly differ from that for weak beams. The main difference stems from the beam energy density, $n_b \varepsilon_b$, being much greater than the gas kinetic pressure of the ambient plasma, $n_0 T_0$. This makes the BPI/BPD regime at the start of the breakdown drastically differ from the presented above (Volokitin and Mishin, 1978; Mishin and Ruzhin, 1980a, 1981).

5.2. Initial Stage of the BPD

5.2.1. Beam Instability

Following Alekhin et al. (1972) and Gendrin (1974), let us consider a cylindrical (“pencil”) beam injected with the beam current, $I_b \ll I_{cr}$, divergence angle, $\Delta\theta_0 < \theta_* = (\omega_{ce}/\omega_{p0}) (I_b/I_{lim})^{1/2} \ll 1$, and initial radius, r_0 . Here the limiting current, $I_{lim} = \varepsilon_b v_b / e \approx 30 \cdot \tilde{\varepsilon}_b^{3/2} [\text{A}]$, with $\tilde{\varepsilon}_b = 0.1 \varepsilon_b [\text{keV}]$, defines the injection current at which the beam is locked by the space charge. The initial beam density, $n_b^{(0)} = I_b / e u_b \pi r_0^2$, greatly exceeds $n_0 \sim 10^5 \text{ cm}^{-3}$ for currents $I_b > 0.1 \text{ A}$ and $\varepsilon_b < 10 \text{ keV}$ ($\tilde{\varepsilon}_b < 1$). Therefore, the beam expands radially due to electrostatic repulsion. For $\omega_{pe}(n_0) = \omega_{p0} > \omega_{ce}$ and injections at small pitch angles $\theta_0 < \theta_*$, the beam quickly (in $\sim \omega_{p0}^{-1}$) expands to

$$r_{b*} \approx (\nu_b / \omega_{p0}) (I_b / I_{lim})^{1/2} \approx 50 \tilde{\varepsilon}_b^{-1/4} \sqrt{I_b / \tilde{n}_0} \text{ cm}, \quad (35)$$

so its density reduces to $n_{b*} \approx n_0 = 10^5 \cdot \tilde{n}_0 \text{ cm}^{-3}$ and the velocity/pitch-angle scatter increases to $\Delta v_b^* / \nu_b \sim \Delta\theta_* \sim (I_b / I_{lim})^{1/2}$. For $\omega_{p0} < \omega_{ce}$, the beam expands to $r_{b*} (\omega_{p0} / \omega_{ce})$ and its density becomes $n_0 (\omega_{ce} / \omega_{p0})^2$ (Gendrin, 1974).

At injection angles $\theta_0 > \theta_*$, the beam executes a Larmor spiral with a hollow radial cross-section bounded by the beam gyroradius, $r_{cb} = \nu_b \sin \theta_0 / \omega_{ce}$ (cf. Winckler, 1992, **Figure 2**). As for small angles, electrostatic repulsion makes the radial beam thickness of $\delta r_* \approx r_{b*}^2 / (2r_{cb} \cos \theta_0)$, so its density nears to n_0 . A helical structure rapidly transforms into a pencil-like (cf. Winckler, 1992, **Figure 2**) as the excited EB/ECH oscillations scatter beam electrons, thus broadening their pitch-angle distribution and enhancing radial diffusion (see shortly).

A generic form of the distribution function of a cylindrical beam injected at $\theta_0 > \theta_*$ can be represented as

$$F_b(\mathbf{v}, r) = n_b(r) F_{\parallel} \left(\frac{u - u_b}{\Delta u_0} \right) F_{\perp} \left(\frac{v_{\perp} - v_{b\perp}}{\Delta v_{\perp 0}} \right) \quad (36)$$

with the beam density, $n_b = I_b / (\pi e r_{\perp}^2 u_b)$, and the radial thickness, $r_{\perp}$. Here $u_b = \nu_b \cos \theta_0$, $v_{b\perp} = \nu_b \sin \theta_0$, $\Delta u_0 \ll u_b$, $\Delta v_{\perp 0} \ll v_{b\perp}$, $\int F_{\parallel} F_{\perp} d^3 v = 1$, and $F_{\parallel, \perp}(x)$ has a maximum (“bump”) at $x \rightarrow 0$ and tends to 0 at $|x| \rightarrow \infty$. Taking $r_{\perp} \sim r_{cb} = \nu_b \sin \theta_0 / \omega_{ce}$, gives $n_b \approx 10^3 I_b \tilde{\varepsilon}_b^{-3/2} / \sin^2 \theta_0 \cos \theta_0 \text{ cm}^{-3}$. Using a hollow-cylinder beam does not change the basic results concerning the instability development.

In a uniform beam-plasma system, a “cold,” $\Delta u/u_b < (n_b/n_e)^{1/3} < 1$, beam excites Langmuir oscillations, $\omega_{k_0} \approx \omega_{p0} \approx k_0 u_b$, at the growth rate $\gamma_c \approx \omega_{p0} (n_b/n_0)^{1/3}$ (e.g., Mikhailovskii, 1974). For a bounded beam of a radial extent $r_\perp$, the growth rate depends on whether the inhomogeneity parameter, $\xi_0 = k_0 r_\perp$, is smaller or greater than the root of the first kind, first order Bessel function, $\xi_1 \approx 3.75$. That is, short-scale waves, $\xi_0 \gg 1$, develop as for an unbounded beam, in accordance with general considerations. For $\theta_0 < \theta_*$, the radius $r_\perp$ is $\approx r_{b*}$ (35) and $\xi_0 \approx (I_b/I_{\text{lim}})^{1/2} \ll 1$. For a “tenuous” beam, $n_b < n_0 \omega_{ce}^2/(\omega_{p0}^2 - \omega_{ce}^2)$, the growth rate becomes (Alekhin et al., 1972)

$$\gamma_c^{(<)} \approx \omega_{p0} (I_b/2I_{\text{lim}}\xi_1^2)^{1/3} \quad (37)$$

For a “dense” beam, $n_b > n_0 \omega_{ce}^2/(\omega_{p0}^2 - \omega_{ce}^2)$, waves at $\omega_k \approx \sqrt{\omega_{p0}^2 - \omega_{ce}^2} \approx k u_b$ grow at a rate

$$\gamma_c^{(>)} \approx \omega_k \left(\frac{\omega_{ce}}{\omega_{p0}} \right)^{1/2} \left(\frac{I_b}{I_{\text{lim}}} \right)^{1/4} \quad (38)$$

The radial extent of the wave excitation region, and thus of the discharge,

$$R_d \sim \frac{v_b}{\omega_{ce}} \begin{cases} |\sin \theta_0| (\xi_1/\xi_0)^{1/3} & \text{at } \theta_0 > \theta_* \\ (I_{\text{lim}}/I_b)^{1/4} & \text{at } \theta_0 \leq \theta_* \end{cases} \quad (39)$$

significantly exceeds the beam gyroradius, particularly, at small injection pitch angles. In fact, the radial size is even greater because the wave amplitudes outside R_d (39) decrease as $\sim (R_d/r)^{1/2} \exp(-r/R_d)$ (Alekhin et al., 1972). This agrees with the laboratory (e.g., Fainberg, 1962; Kharchenko et al., 1962; Bernstein et al., 1978; Jost et al., 1982) and mother-daughter (Maehlum et al., 1980b,c; Grandal, 1982b; Jacobsen, 1982; Duprat et al., 1983) measurements.

As already noted, the instability due to the bump in transverse velocities is important at $\theta_0 \gg \theta_*$. At $\Delta v_{\perp 0} \rightarrow 0$, the distribution function, $F_\perp \left(\frac{v_\perp - v_{b\perp}}{\Delta v_{\perp 0}} \right)$ tends to the Dirac delta function, $\delta(v_\perp - v_{b\perp})$, known and as a “ring” or “oscillator” distribution (e.g., Mikhailovskii, 1974). In a uniform beam-plasma system, EB/ECH oscillations can develop with almost any possible ratio of $\xi_\perp = k_\perp v_{b\perp}/\omega_{ce}$. Similar to the bump-in-tail instability, the beam radial inhomogeneity is less significant for short-scale, $\xi_\perp \gg 1$, oscillations. For those, the maximum growth rate is reached at $\omega_k \approx \omega_{uh} \approx k_z u_b + s\omega_{ce}$

$$\gamma_o \approx \omega_{uh} \left(\frac{n_b}{\pi n_0} \right)^{1/3} \left(\frac{k_z}{k} \right)^{2/3} \xi_0^{-1/3} \quad (40)$$

If $k_z \rightarrow 0$, the maximum growth rate at $\omega_k \approx \omega_{uh} \sim s\omega_{ce}$ ($s \geq 2$) reduces to $\gamma_s \approx 0.1 \left(\omega_{p0}^2/\omega_{ce} \right) (n_b/n_0)^{1/2}$ and at $s \gg 1$ tends to $\gamma_s \approx \omega_{p0} (n_b/n_0)^{1/2}$, where $a_s = 1/2$ at $n_b/n_0 < \omega_{ce}^2/\omega_{p0}^2$ and $2/5$ otherwise (e.g., Mikhailovskii, 1974).

5.2.2. Relaxation of Cold Beams

As the excited spectrum is narrow, $\Delta k \ll \gamma_c/u_b$, and width of the resonance region, $|u - v_{ph}| \sim \gamma_c/k_0$, is greater than Δu_0 , it is safe considering interaction with a quasi-monochromatic Langmuir oscillation. Henceforth, γ_c stands for $\gamma_c^{(<)}$ or $\gamma_c^{(>)}$, whichever applicable. The wave growth slows down when the beam electrons become trapped by the wave potential and change the relative phase bouncing back and forth in the potential hole. The beam velocity scatter, $\Delta u_{tr}/u_b$, increases over the saturation time $\sim 10/\gamma_c$ to $\sim \gamma_c/\omega_{p0}$ at the wave amplitude (e.g., Onishchenko et al., 1970; van Wakeren and Hopman, 1972; Abe et al., 1979)

$$|E_0^{(tr)}| \approx (8\pi n_b \varepsilon_b \gamma_c/\omega_{p0})^{1/2} |\cos \theta_0| \quad (41)$$

That is, the instability is saturated along the relaxation length, $l_{rel}^{(tr)} \sim 10u_b/\gamma_c \sim 100\text{--}300$ m.

The saturation of the ring/oscillator instability also occurs due to trapping of the beam electrons at the wave amplitude $|E_\theta^{(tr)}| \approx |E_\parallel^{(tr)}|$ (41) with γ_c and $\cos \theta_0$ replaced by γ_{osc} or γ_s and $\sin \theta_0$, respectively (e.g., Kitsenko et al., 1974; Aburjania et al., 1978). In the saturated state, the scatter of perpendicular velocities is of the order of $\gamma_{o,s}/\omega_{uh} \sim 0.2\text{--}0.3$. Since the wave frequency exceeds the cyclotron frequency, the trapped beam electrons are unmagnetized and pulled by the wave across the magnetic field, ultimately filling the void in the center and expanding over the beam gyroradius. This is consistent with the data concerning the beam structure (Bernstein et al., 1979), as well as the distortion of single particle trajectories for beam currents above the threshold (Maehlum et al., 1980c).

5.2.3. Electron Heating and Ionization

If the beam is relatively weak and $W_{tr} = |E_0^{(tr)}|^2/4\pi < n_0 T_0$, the electron heating rate is given by equation (27) with the coefficient of inelastic losses, $\delta_{il}(T_h)$, adjusted to higher values of $T_h > 1$ eV. However, $W_{tr} > n_0 T_0$ for $I_b \geq 0.1$ A and $\varepsilon_b \leq 10$ keV, so the quiver velocity, v_E (10), for the wave (41) is greater than the thermal electron velocity and $r_E = v_E/\omega_{p0} > r_D(T_0)$. This results in excitation of secondary Langmuir waves, E_p , and low-frequency, $\Omega \leq \mu^{1/2}\omega_{p0}$, density oscillations, δn_p , via an aperiodic parametric instability, $E_0^{(tr)} \rightarrow E_p + \delta n_p$, with the maximum growth rate, $\gamma_p \sim \mu^{1/3}\omega_{p0}$, at $\mathbf{k} = \mathbf{k}_p \approx 1.8\mathbf{r}_E/r_E^2$ (Kruer and Dawson, 1971; De Groot and Katz, 1973). Trapping of thermal electrons by the secondary waves makes the electron orbits intersect at $\tau_p \approx 10\gamma_p^{-1} \sim 20$ μ s. This leads to the fast, less than f_{p0}^{-1} , heating of the bulk electrons up to $T_h \approx T_E \approx \frac{1}{2} m_e v_E^2$ (De Groot and Katz, 1973). Substituting $|E_0^{(tr)}|$ (41) into v_E (10), yields

$$T_h = \varepsilon_{ion} \cdot \tilde{T}_h \approx \varepsilon_b \frac{n_b}{n_0} \frac{\gamma_c}{\omega_{p0}} \cos^2 \theta_0 \quad (42)$$

That is, the collisionless heating time, $\tau_{heat}^{(c)} \sim 10(1 + \gamma_c/\gamma_p)/\gamma_c$, is much faster than the collisional time. For the ring/oscillator

beam, the value of T_h follows from $|E_{\theta}^{(tr)}|$. It is worth noting that the fast heating was observed in the Bauer et al. (1992) laboratory experiment. High electron temperatures in the near zone were reported by Gringauz et al. (1981, up to ~ 100 eV), Jacobsen (1982), and Arnoldy et al. (1985).

At $n_0 \sim 10^5 \text{ cm}^{-3}$ and $1.5 < \tilde{T}_h < 10$, the ionization rate, $\nu_{ion}(T_h) \sim 10^{-8} \tilde{T}_h^{1/2} N$, exceeds τ_p^{-1} at

$$N > N_{hf} \approx 10^{10} \sqrt{n_0 / \tilde{T}_h} \sim 3 \cdot 10^{12} / \tilde{T}_h^{1/2} \text{ cm}^{-3} \quad (43)$$

or $h < h_{hf} \approx 110 - 120$ km. Here the BPD criterion is alike the HF discharge (32), while at higher altitudes it becomes $\nu_{ion}^{-1}(T_h) < \tau_{loss}$. The lifetime of the heated electrons is determined mainly by transverse diffusion, $\tau_{d\perp} \sim R_d^2 / \nu_e(T_h) r_{ce}^2$. However, the time of injection into the same magnetic tube is limited due to the transverse speed of the beam guiding center, $\tau_R \sim r_{cb} / V_{\perp}$, with $\mathbf{V}_{\perp} = \mathbf{V}_{R\perp} + \Delta \mathbf{V}_{\perp}(\psi)$ (see section 3.3). It is worth noting that Vlasenko et al. (1976a) have observed that the BPD is inhibited at some critical speed of the gas flow across the electron beam.

At high altitudes, the BPD condition reduces to $\tau_R > \nu_{ion}^{-1}(T_h)$, which yields (Mishin and Ruzhin, 1980a)

$$N > N_{\min} \approx 10^{10} V_{R\perp} \begin{cases} 3\tilde{\varepsilon}_b^{1/4} I_b^{-1/2} & \text{at } \theta_0 \leq \theta_* \\ \frac{\omega_{p0}}{2\omega_{ce}} I_b^{-2/3} & \text{at } \theta_0 > \theta_* \end{cases} \text{ cm}^{-3} \quad (44)$$

For injection currents ~ 0.5 A, $\tilde{\varepsilon}_b \sim 1$, and $V_{R\perp} \sim 0.3$ km/s, the BPD condition (44) yields $h(N_{\min}) = h_{\max} \sim 160 - 170$ km.

A few remarks are in order. Interaction of beam electrons with short-scale, $k \gg \omega_{pe} / v_b$, oscillations results in fast pitch-angle scattering with the effective “collision” frequency (Mishin et al., 1989; Mishin et al., 1994; Khazanov et al., 1993)

$$\nu_{eff} \approx \omega_{pe} \sum_k \frac{W_k}{kr_D n_e T_h} \left(\frac{T_h}{\varepsilon_b} \right)^{3/2} \quad (45)$$

which greatly exceeds the collisional frequency, $\sim \nu_b$, at altitudes above ~ 120 km. Numerical simulations (Khazanov et al., 1993; Mishin et al., 1994) have shown that this process suffices to explain the basic features of the prompt electron echo (section 3.4).

Thus far, only perpendicular diffusion was considered. However, at injection angles close to $\pi/2$ the field-aligned extent of the discharge gap, $\sim l_{rel}^{(tr)} \cdot \theta_*$, is so small that the parallel diffusion becomes dominant and limits the discharge ignition. Therefore, besides parallel injections, the optimal BPD conditions are achieved at injection angles $\leq 80^\circ$ (cf. Figure 5).

Next, the conversion rate (26) increases to $\gamma_{conv} \approx \omega_{pe} (|\delta n_k|^2)$ for $k \sim r_D^{-1}$ (Volokitin and Mishin, 1978). In the heated plasma, the low-frequency density oscillations, $k \sim k_p \sim r_D^{-1}(T_h)$, are saturated at $|\delta n_p|^2 \sim 2\mu^{1/3}$ (De Groot and Katz, 1973), thus yielding $\gamma_{conv}^{(p)} \approx 2\mu^{1/3} \omega_{pe}$. If $\gamma_{conv}^{(p)} > \gamma_c$ ($\gamma_{conv}^{(p)} > \gamma_{osc}$ or γ_s), the beam (ring) instability is suppressed. That is, the heated plasma is “cleared” for the beam propagation and the BPI starts in the adjacent region. This way, the heated region will move

away from the rocket at a speed $u_T^{(p)} \sim l_{rel}^{(tr)} / \tau_p$ (Mishin and Ruzhin, 1980a, 1981).

The short-scale density oscillations in the “cleared” region decay due to various dissipative processes including diffusion and induced scattering on ions within the decay time $\tau_{osc} \sim 10\tau_{heat}^{(c)}$. As soon as $\gamma_{conv}^{(p)}$ drops below γ_c (γ_o or γ_s), the BPI resumes in the initial region and is suppressed again after the heating time, $\tau_{heat}^{(c)}$. The obvious corollary is that the BPI in each individual gap, $\sim l_{rel}^{(tr)}$, proceeds in a quasi-periodic series of short, $\sim \tau_{heat}^{(c)}$, pulses of enhanced wave activity and particle energization repeated at $\sim \tau_{osc}$. Such behavior is typical of laboratory experiments with intense cold beams (e.g., Kharchenko et al., 1962; Cabral, 1976; Vlasenko et al., 1976b) and has also been observed in active experiments (e.g., Gringauz et al., 1981; Kawashima, 1988). Evidently, at altitudes above $h_{\max}$ this “oscillatory” regime persists over the duration of injection pulses.

During the oscillation period, τ_{osc} , the newly-born plasma cools off due to thermal conduction and inelastic losses down to $T_e \sim 1$ eV and the ionization rate decreases. However, as the diffusion time is greater than τ_{osc} , the plasma density, n_e , in the heated region increases with the average rate $\bar{\tau}_{ion}^{-1}$. As long as $n_b^{(0)}$ remains greater than n_e , the “initial stage” regime holds. When n_e rises to $n_b^{(0)}$ at the time $\tau_0 \sim \bar{\tau}_{ion} \ln(n_b^{(0)} / n_0) \sim 10N_{\min} / N$ ms, electrostatic repulsion ends and the beam preserves its initial snape. However, the instability development proceeds as before and even faster because of the increase of ω_{pe} . After a few τ_0 , the plasma density, n_e , significantly exceeds $n_b^{(0)}$ so that the “warm”-beam approximation (section 4) becomes applicable. Actually, that can happen even earlier due to radial diffusion of beam electrons scattered off persisting short-scale oscillations at the rate (45), which is indicated by the prompt electron echo (section 3.4).

5.3. Stationary BPD

There are the principal differences between the stationary and initial BPD regimes. First, a cold beam excites convective modes with the group velocity $v_g \sim u_b$, while a warm beam excites Langmuir waves (6) with $v_g \ll u_b$. As result, the energy density of the excited waves remains very high even for $n_b \ll n_e$, as well as the electron temperature. Next, in addition to ionization by the heated thermal plasma electrons, accelerated suprathermal electrons can contribute significantly to the BPD ignition when strong Langmuir turbulence determines the beam relaxation (Mishin and Ruzhin, 1980a, 1981; Rowland et al., 1981a; Papadopoulos, 1982, 1986; Sharp, 1982; Omelchenko et al., 1992; Sotnikov et al., 1992). If the SLT development is inhibited, the bulk electron heating is the only source for the BPD. Here the BPD development is similar to that of the HF discharge, as discussed in section 5.1.

As the neutral density in the NRG region varies during the flight, the beam relaxation regime changes accordingly. Let us first consider the region where $n_b > n_b^{(c)}$ (22). At $v_e < v_*$ (23), as follows from the heat balance (31) with $\gamma_b / \omega_{pe} \sim 10^{-2} - 10^{-1}$, the electron temperature is $T_h \sim 1 - 5$ eV. The ionization by the thermal bulk electrons is determined by a

Maxwellian tail, i.e., $q_M \sim 10^{-8} n_e N \exp(-\tilde{T}_h^{-1})$. However, the main contribution comes from suprathermal electrons (29). Comparing the ionization rate $q_{ion}^{(t)}$ (30) with τ_R yields the maximum plasma density in a stationary discharge of the order of

$$n_{\max}^{(t)} \sim n_b v_e (T_h) \frac{10 T_h}{\varepsilon_{ion}} \left(\frac{u_b}{\Delta u_b} \right)^4 \frac{r_{\perp}}{V_{R\perp}} \propto \frac{I_b}{u_b} \frac{N}{V_{R\perp}} \quad (46)$$

This dependence is similar to the observed $f_{\max} \propto \sqrt{n_{\max}}$ (4).

As soon as $v_e > v_*^*$ or $N > N_{md}$, the electron temperature increases but the acceleration of suprathermal electrons is inhibited. Here the plasma density is determined by the balance between ionization by thermal electrons and radial diffusion to give $n_{\max}^{(t)} \propto N^{2/5}$ and then $\propto N^{3/2}$, as collisional damping reduces the beam spreading and hence the excited wave energy. The lower BPD boundary is determined by the neutral density, $N_{\max}$, at which collisional damping inhibits the development of the cold-beam instability (Mishin and Ruzhin, 1980a; Dokukin et al., 1981).

These results pertain to the beam instability. As far as the ring/oscillatory instability is concerned, its development at $\omega_{pe}/\omega_{ce} \sim s \gg 1$ proceeds even faster than at the initial stage. As a result the beam distribution over transverse velocities widens such that the wave energy grows to $\sim (0.1 - 0.3) n_b \varepsilon_b$. Here, the strong electron heating is the main contributor to the BPD development. Since the group velocity at $\omega_k/\omega_{ce} \approx s + 1/2$ tends to zero, these waves are most enhanced (cf. Ashour-Abdalla and Kennel, 1978; Ashour-Abdalla et al., 1980), which explains the gyrofeatures in **Figure 7B**. As the ring instability growth rates are smaller than γ_c , it is suppressed at higher altitudes. **Figure 12** illustrates the BPD regimes vs. altitude and the range of altitudes (neutral gas densities) where the various BPD features were observed.

Finally, using the heated bulk and accelerated tail electrons it is easy to show that the total power of optical emissions radiated from the NRG is of the order of a few per cent of the beam power (Ivchenko et al., 1981; Mishin and Ruzhin, 1981). The power of VHF radioemission estimated assuming conversion of Langmuir waves on density oscillations inherent in unstable beam-plasma systems also reaches a few percent of the beam power (Galeev et al., 1976; Mishin and Ruzhin, 1981).

6. ENHANCED AURORA

The distribution of precipitating electrons, $F_b(\mathbf{v})$, varies significantly in time and space creating various auroral forms with latitudinal scale lengths from tens of kilometers down to hundred meters (e.g., Meng et al., 1991). Some of the measured distributions of auroral electrons do exhibit the bump-in-the-tail feature. Its origin is beyond the scope of this survey.

6.1. Auroral Electrons

Figure 13 exemplifies auroral beam and suprathermal electron spectra observed over auroral arcs (e.g., Reasoner and Chappell, 1973; Arnoldy et al., 1974; Feldman and Doering, 1975; Bryant et al., 1978). In **Figure 13B**, the primary flux is approximated by a Gaussian distribution of the density $n_b = 0.6 \text{ cm}^{-3}$,

energy scatter $\Delta \varepsilon_b \approx 2.9 \text{ keV}$, and $\varepsilon_b \approx 10.2 \text{ keV}$ (the dashed line). These beam parameters easily suffice the BPI conditions described in sections 4.3 and 4.4. Therefore, the beam relaxation can be described by the BPI theory outlined there. In particular, the relaxation length, l_b (25) at $v_e < v_*$ (23) exceeds the distance between the acceleration region of auroral beams (e.g., Meng et al., 1991) and the E-region ionosphere. This explains why the bump-in-tail distributions, such as in **Figures 13A,B**, are preserved along the path (Galeev, 1975; Papadopoulos, 1975).

As in **Figures 6A,B**, the suprathermal spectrum in the range $\geq 6\text{--}1,000 \text{ eV}$ is well approximated by a power law $\Phi(\varepsilon) \sim \varepsilon^{-1}$. It is considerably flatter than the SPA spectrum (3) shown by the solid line in **Figure 6B**. In frame c, a SPA-like spectrum over the class II arc changes to a flatter one at $\varepsilon \geq \varepsilon_{\min} \approx 20 \text{ eV}$, i.e., $\varepsilon_{\min} \gg \varepsilon_{\min}^{(M)}$. The overall observations show that the spectrum of suprathermal electrons over arcs is not formed solely on account of the collisional interaction but can be explained in terms of the SLT acceleration (Papadopoulos and Coffey, 1974; Galeev, 1975; Matthews et al., 1976; Mishin and Telegin, 1986). It is worth to note that recent incoherent scatter radar observations (Isham et al., 2012; Akbari et al., 2013) do reveal the signatures of strong Langmuir turbulence in auroral plasma, quite similar to that in high-power radio wave experiments.

6.2. Luminosity and Ionization Profiles

A typical representation of the altitude-profile of auroral luminosity/ionization is illustrated by **Figure 2**. However, Donahue et al. (1968) reported on rocket measurements of auroral emissions at 557.7 and 391.4 nm, both having narrow local maxima at 115 and 130 km. Similarly, ground-based optical imagers detected either double-peaked auroral rays of about the same thickness, displaced in altitude by about 5–15 km, or one thick layer with the sharp upper boundary below about 130 km, with the characteristic scale length of only a fraction of H_N (Oguti, 1975; Stenbaek-Nielsen and Hallinan, 1979; Dzyubenko et al., 1980; Hallinan et al., 1985). This phenomenon was called “Enhanced Aurora.” Similar layers of auroral ionization were also detected from sounding rockets (Swider and Narcisi, 1977; Morioka et al., 1988), the EISCAT UHF incoherent scatter radar (Wahlund et al., 1989; Schlesier et al., 1997), and a dual-altitude 90-MHz radar system (Timofeev and Miroshnikov, 1982). There were efforts to explain such double-peak profiles by the collisional interaction invoking two precipitating electron populations, i.e., a monoenergetic field-aligned beam and an isotropic beam of a higher energy. However, these efforts failed to fit the upper peak and sharp upper boundary.

Figures 14A,B shows double-peaked auroral rays observed near Tixie Bay by a side-looking low-light TV camera, the same as used in Zarnitza 2, and their luminosity profiles (Dzyubenko et al., 1980). The upper peaks are by a factor of two narrower than the minimum possible from the SPA, while the lower peak matches the SPA predictions (cf. **Figure 4**). Shown below (**Figure 14C**) is an example of a rayed arc with the sharp upper boundary (Hallinan et al., 1985).

More than fifty narrow layers of the enhanced electron temperature co-located with the plasma density peaks in

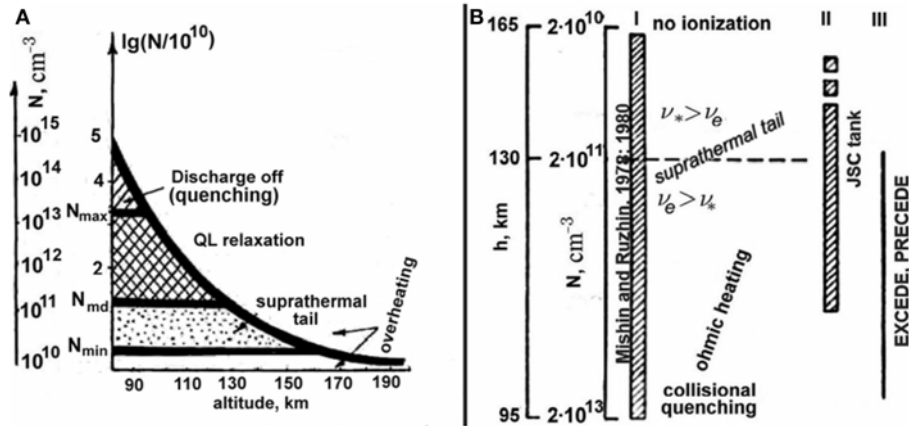


FIGURE 12 | (A) The different stationary BPD regimes vs. altitude defined by the beam relaxation regimes. Adapted from Mishin and Ruzhin (1980a). **(B)** Altitude and neutral density range over which BPD is expected. The first stripped column indicates the range expected from Mishin and Ruzhin's analysis. The second stripped column indicates the parameter range over which BPD has been observed in the large tank at the Jonson Space Center. The vertical line on the right indicates the altitude range over which the EXCEDE and PRECEDE experiments with currents far beyond the BPD threshold were conducted. Adapted from Linson (1982). Reprinted by permission from Plenum Press.

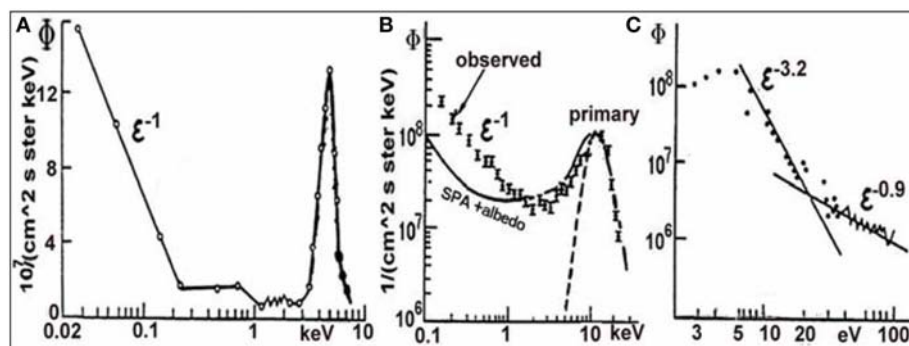


FIGURE 13 | Differential number fluxes over auroral arcs with the 'bump-in-the tail' distribution of primary electrons and flat power-law spectra of secondary electrons. Dots and circles show the observed fluxes, while dashed and solid lines show (B) a Gaussian approximation and (C) power-law trends. Adapted from (A) Arnoldy et al. (1974). (B) Bryant et al. (1978), and (C) Feldman and Doering (1975). Reprinted by permission from the American Geophysical Union.

the altitude range ≥ 115 –150 km have been found in the EISCAT UHF radar database (Schlesier et al., 1997). One sample is shown in **Figures 15A,B**, where the thin layers in T_e and n_e are emphasized by thick lines. Note that the ion temperature is significantly smaller than the peak value of $T_e \approx 3,000$ K and that the overall density profile has two peaks. A similar double-peaked ionization profile (**Figure 15C**) was observed in pulsating aurora (Wahlund et al., 1989). Dashed lines show the SPA profiles calculated for $\epsilon_b = 3.8$ and 10 keV. As for auroral rays, the difference with the SPA profile for the upper peak is evident.

A striking resemblance between the Enhanced and Artificial Aurora (**Figure 4**) profiles is evident. Given that the parameters of electron beams far beneath the rocket are close to that of natural beams, it is safe to conclude that their generation mechanisms have much in common. It is obvious that the sought-for mechanism is the one that creates the plasma turbulence layer.

7. CONCLUSION

The effects of powerful electron beams injected from sounding rockets into the upper atmosphere to create artificial aurora are outlined. Data come from *in situ* measurements of the luminescence, thermal and suprathermal populations, and beam electrons near a beam-emitting space vehicle, as well as from ground-based optical, radar, and radioemission observations. The overall dataset cannot be explained solely by collisional degradation of energetic electrons but demands collisionless beam-plasma interactions (BPI) be taken into account. A brief survey of the BPI theory in a weakly-ionized plasma is presented. The basic processes of the near-rocket region are described in terms of the beam-plasma discharge (BPD) ignited by plasma electrons energized by the beam-excited plasma turbulence. Depending on the ambient plasma and atmospheric densities, there are several regimes of the BPD development. The observations of artificial auroral rays far beneath the rocket

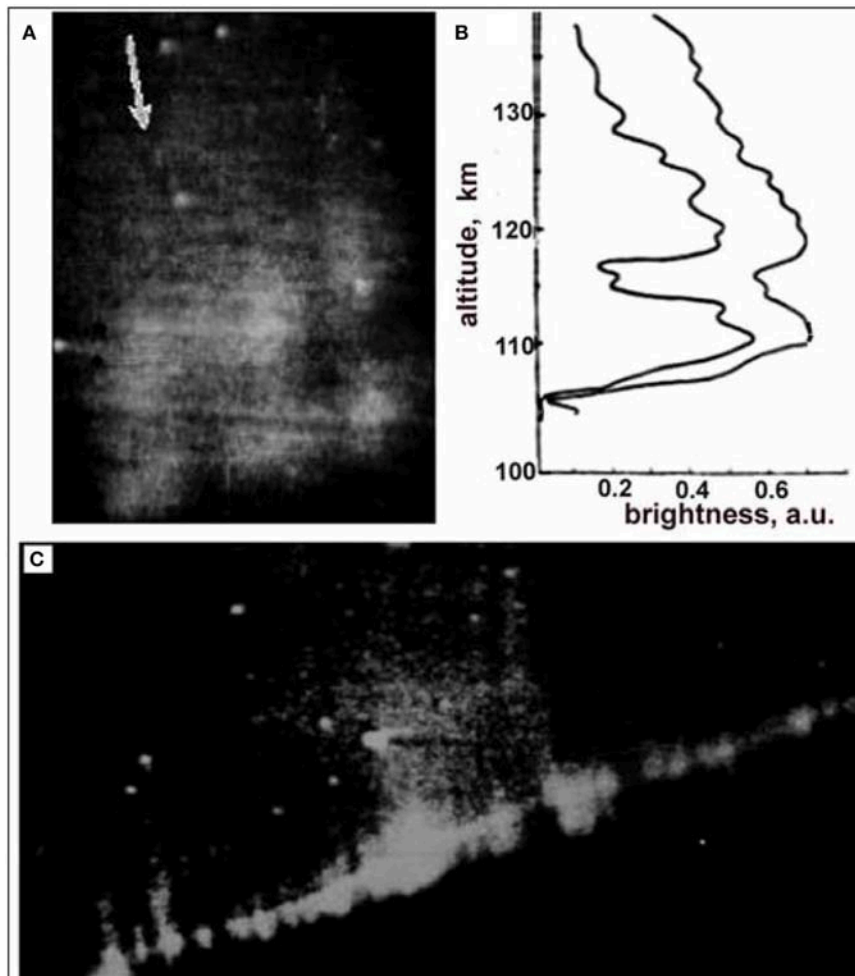


FIGURE 14 | (A) Auroral rays along the magnetic field indicated by the arrow and **(B)** their luminosity profiles. Adapted from Dzyubenko et al. (1980). **(C)** A rayed arc with the sharp upper boundary. After Hallinan et al. (1985). Reprinted by permission from the American Institute of Physics and American Geophysical Union.

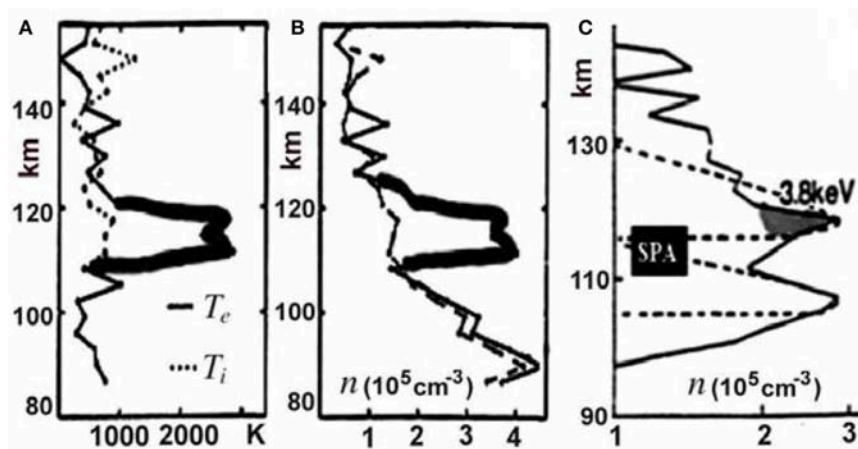


FIGURE 15 | Thin layers of the **(A)** elevated electron temperature and **(B)** electron density indicated by thick lines (after Schlesier et al., 1997); and **(C)** a double-peaked ionization profile in pulsating aurora. Dashed lines show the SPA profiles calculated for $\epsilon_D \approx 3.8$ and 10 keV. After Wahlund et al. (1989). Reprinted by permission from the American Geophysical Union.

indicate that turbulence regime and thus the relaxation of radially expanded beams are strongly affected by collisions of plasma electrons. As a result, the energy density of plasma waves and concomitant energization of plasma electrons are enhanced in a narrow layer termed the plasma turbulence layer (PTL). The PTL formation results in an upper peak in a double-peak structure of artificial auroral rays. Some examples of optical and incoherent scatter radar observations of the luminosity and ionization profiles of rayed auroral arcs exhibiting two peaks or a sharp upper boundary are presented. Such auroral forms have been called Enhanced Aurora. An evident resemblance between Enhanced and Artificial Aurora points to their common

generation mechanism, which is the one that creates the plasma turbulence layer.

AUTHOR CONTRIBUTIONS

The author confirms being the sole contributor of this work and has approved it for publication.

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NOTATION

B_0 - the ambient magnetic field
 L - Langmuir waves
 EB - electron Bernstein waves
 F_M - the Maxwellian distribution
 F_s - the distribution function of secondary electrons
 F_t - the accelerated tail distribution function
 MI - modulational instability
 Primary waves - the waves excited by the beam
 SLT - strong Langmuir turbulence
 SST - superstrong (Langmuir) turbulence
 T_{spin} - the rocket's rotation period
 WT - weak turbulence
 W_k - the wave spectral energy density
 $l_n = d \ln n_e / dh|^{-1}$ - the plasma density scale height
 n_t^{ion} - the density of the ionizing ($\varepsilon > \varepsilon_{ion}$) electrons
 r_D - the Debye radius
 r_{cb} - the Larmor radius of the beam electrons
 $w = W/n_0 T_e$ - the dimensionless parameter of nonlinearity
 w_{th} - the modulational instability threshold
 γ_{mi} - the modulational instability growth rate
 $\alpha_t^{ion} = n_t^{ion}/n_c$ - the relative density of ionizing ($\varepsilon > \varepsilon_{ion}$) electrons
 ε_{ion} - the ionization energy
 $\varepsilon_{min} (\varepsilon_{max})$ - the minimum (maximum) energy of the accelerated electrons
 θ - pitch angle
 λ_T - the mean free path of thermal electrons
 $\mu = m_e/m_i$ - the electron-to-ion mass ratio
 ν_{il} - the frequency of inelastic collisions
 ν_{ion} - the ionization frequency



Mutual Impedance Probe in Collisionless Unmagnetized Plasmas With Suprathermal Electrons—Application to BepiColombo

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Context: Mutual impedance experiments are active electric probes providing *in-situ* space plasma measurements. Such active experiments consist of a set of electric antennas used as transmitter(s) and receiver(s) through which various dielectric properties of the plasma can be probed, giving therefore access to key plasma parameters such as, for instance, the electron density or the electron temperature. Since the beginning of the space exploration, such active probes have been launched and operated in Earth's ionospheric and magnetospheric plasmas. More recently and in the coming years, mutual impedance probes have been and will be operated onboard exploratory planetary missions, such as Rosetta, BepiColombo and JUICE, to probe the cometary plasma of 67P/Churyumov-Gerasimenko, the Hermean and the Jovian magnetospheres, respectively.

Aims: Some analytic modeling is necessary to calibrate and analyse mutual impedance observations in order to access to macroscopic bulk plasma quantities. *In situ* particle observations from various space missions have confirmed that space plasmas are out of local thermodynamic equilibrium. This means that particle velocity distributions can be far from a Maxwellian distribution, exhibiting for instance temperature anisotropies, beams or a suprathermal population. The goal of this paper is to characterize the effect of suprathermal electrons on the instrumental response in order to assess the robustness of plasma diagnostics based on mutual impedance measurements in plasmas characterized by a significant amount of suprathermal particles.

Methods: The instrumental response directly depends on the electron velocity distribution function (evdf). In this work, we choose to model suprathermal electrons by considering different approaches using: (i) a kappa evdf, (ii) a double-Maxwellian evdf or (iii) a mix of a Maxwellian evdf and a kappa evdf. For each case, we compute the spatial distribution of the electrostatic potential induced by the transmitters, discretized and modeled here as an ensemble of pulsating point charges.

Results: We apply our modeling by building synthetic mutual impedance spectra of the PWI/AM²P probe, launched in October 2018 onboard the Mercury Magnetospheric Orbiter (MIO/MMO) spacecraft of the BepiColombo exploratory space mission, in order to calibrate and analyse the future electron observations in the plasma environment of Mercury.

Keywords: mutual impedance experiments, modeling, electrostatic radiated potential, BepiColombo, mercury, suprathermal electrons, active experiment

1. INTRODUCTION

Mutual impedance experiments are active electric experiments designed to measure *in-situ* space plasma bulk properties such as the electron density and the electron temperature (Chasseriaux et al., 1972). The measurement is usually based on the electric coupling between pairs of electric dipole antennas embedded in the plasma to be probed (Storey et al., 1969). The transmitting electrodes inject an oscillating current at a given frequency in the surrounding plasma. This current and the electric potential difference induced on the receiving antenna are both measured simultaneously at the same frequency. A mutual impedance spectrum is built by varying, step by step, the emitted frequency.

Initially developed in geophysical fields prospecting to measure the resistivity of the ground (Wenner, 1915; Schlumberger, 1920; Storey et al., 1969), mutual impedance experiments have been used on Earth ionospheric and magnetospheric missions (Beghin and Debie, 1972; Décréau et al., 1978; Beghin et al., 1982). More recently, mutual impedance experiments have been used to probe interplanetary plasmas. The Mutual Impedance Probe (MIP), as a part of the Rosetta Plasma Consortium (RPC), on board the Rosetta orbiter (Trotignon et al., 2007), measured the electron density in the ionosphere of the comet 67P/Churyumov-Gerasimenko (Henri et al., 2017). The Active Measurement of Mercury's Plasma (AM²P) instrument (Trotignon et al., 2006) from the Plasma Wave Investigation (PWI) is currently onboard the Mercury Magnetospheric Orbiter (MIO/MMO) of the BepiColombo mission successfully launched in October 2018. After the 7.2 years cruise phase, this experiment will constrain the plasma bulk properties in the Hermean magnetosphere. Two others experiments from the PWI consortium will operate in the Hermean magnetosphere and in the solar wind close to Mercury in order to measure the electron density onboard the MIO/MMO spacecraft. First, a thermal electrostatic noise spectroscopy experiment (PWI/SORBET, Moncuquet et al., 2006) will operate using the WPS antenna (Benkhoff et al., 2010; Kasaba et al., 2010). This experiment makes use of passive measurements combined to the Quasi-Thermal Noise spectroscopy technic to access the plasma bulk properties such as the electron density or the electron temperature through a diagnostic of the voltage power spectrum (Meyer-Vernet et al., 2017). Second, the spherical probes located at the end of the two MEFISTO antennas will be operated using the Langmuir Probe measurement technique to also access the plasma bulk properties (Blomberg et al., 2006). A strong advantage of the MIO/MMO

spacecraft of BepiColombo is that it is the first time a single spacecraft will carry these three experiments that will be operated simultaneously to provide bulk plasma measurements, thus enabling to take advantage of the strength of each measurement technic and going beyond the intrinsic limitations of each. In the future, the Mutual Impedance MEasurement (MIME) as a part of the Radio Wave Plasma Investigation (RPWI) is being developed for the Jupiter ICy Moons Explorer (JUICE) mission to constrain the Jovian magnetospheric plasma and the ionosphere of Ganymede.

The mutual impedance between two electric antennas immersed in a plasma strongly depends on the plasma properties, in particular the electron velocity distribution function (evdf). As mutual impedance experiments have been used in several plasma environments, many theoretical works have been carried out (Grard, 1969; Navet et al., 1971; Rooy et al., 1972; Pottellette et al., 1975; Beghin, 1995) to characterize the properties of mutual impedance experimental behavior from cold (modeled by a Dirac evdf) to hot (modeled by a Cauchy or Maxwellian evdf) plasmas. However, the impact of high-energy electron called suprathermal electrons, omnipresent in space plasmas, had not been sufficiently considered in the past. The goal of this paper is therefore to fill this gap and study the effect of suprathermal electrons on the instrumental response of mutual impedance experiments.

Indeed, suprathermal electrons are ubiquitous in collisionless space plasmas: in the solar wind (Vasyliunas, 1968), in the Hermean magnetosphere (Christon, 1987; Ho et al., 2016), in the magnetosphere of Saturn (Schippers et al., 2008) or in the ionosphere of the comet 67P/Churyumov-Gerasimenko (Clark et al., 2015; Broiles et al., 2016; Myllys et al., private communication). The evdf in the presence of a suprathermal tail is usually described as the sum of thermal (core) and a non-thermal (halo) parts (Maksimovic et al., 2005):

$$f = f_{core} + f_{halo} \quad (1)$$

where the thermal part is usually described by a Maxwellian evdf and the non-thermal part by a kappa evdf (Lazar et al., 2017). The kappa evdf can be seen as a generalization of the Maxwellian evdf, nearly Maxwellian at low energies and decreases as a power-law at higher energies (Summers and Thorne, 1991). In the literature, observed electron distribution functions have also been modeled by other evdf or combinations of evdf: Maksimovic et al. (1997) fitted the evdf observed by Ulysses in the solar wind with a single kappa evdf, while Schippers et al. (2008), Broiles et al. (2016) and Myllys et al. (submitted) used two kappa evdf to fit the

observed evdf respectively in the Saturn magnetosphere and in the ionosphere of 67P/Churyumov-Gerasimenko.

Previous works investigated the instrumental response of mutual impedance experiments in a presence of suprathermal particles but only in restrictive cases: (i) in a plasma described by monoenergetic evdf [Dirac delta evdf, Grard (1997)] or (ii) in a plasma described by a sum of two Maxwellian evdf on a restricted hot-to-cold electron density and in the limit where the Debye length λ_D is very small compared to the distance between the transmitter and the receiver antennas (Pottelette and Storey, 1981). Recently, Gilet et al. (2017) developed a model of the electrostatic radiated potential in a plasma described by a sum of two Maxwellian evdf down to conditions encountered in interplanetary and planetary plasmas (i.e., $\lambda_D \sim$ transmitter-receiver distance). In this present work, we consider suprathermal electrons associated to a collisionless plasma, for which the hypothesis of thermodynamic equilibrium is no longer valid. In other words, this means that suprathermal particles cannot be considered as a Maxwellian distribution. Instead, we will make use of kappa distributions to model out-of-thermodynamic equilibrium evdf for suprathermal electrons. Especially, we study the robustness of the plasma density measurement through the mutual impedance method in the presence of energetic electrons. This new model is applied to the mutual impedance experiment PWI/AM²P onboard the Mercury Magnetospheric Orbiter (MIO/MMO) of the BepiColombo mission (Trotignon et al., 2006; Benkhoff et al., 2010) to prepare the future calibration of the experiment.

This paper is organized as follows: in section 2, we remind the definition of the electric potential induced by a pulsating point charge in a plasma, when evdf is a combination of kappa and Maxwellian evdf. As mutual impedance experiments are based on the propagation of an electric field in a plasma, we also remind the dispersion relations of the linear eigenmodes of interest of such experiments in section 3. This is done for each considered evdf and it allows to better understand, at least qualitatively, the damping rate of the radiated electric potential in the frequency range encompassing the electron plasma frequency. The electric potentials are then computed and compared to the results obtained from the different evdf such as those are considered in this work. We apply the developed computation to the active quadrupolar mutual impedance probe PWI/AM²P onboard the MIO/MMO spacecraft in section 4. We show that in certain limit (high electron density, small Debye length), the presence of the suprathermal electrons do not change the instrumental response. However, for small enough electron density and large enough Debye length, the more suprathermal electrons are presents, the easier the electron plasma frequency can be measured. This seemingly counterintuitive result is due to the fact that the Debye length is smaller for kappa evdf at equivalent (Maxwellian) temperature. In section 5, we compute the AM²P spectra in typical solar wind plasma and in Hermean magnetospheric plasma, using respectively modeling of evdf from several solar space missions (Pierrard et al., 2016) and the *in-situ* particles measurement from a Mercury flyby by Mariner 10 (Baker et al., 1986). We show how the measurement of the plasma density is not influenced by suprathermal electrons

in typical solar wind plasma close to the Mercury perihelion (0.31 AU) but can be slightly affected close to the aphelion (0.47 AU). Moreover, we show that the detection of the plasma frequency might be challenging in the low density Hermean magnetospheric plasma. Finally we conclude our study in section 6.

2. MODEL

The electric potential ϕ induced in an isotropic, homogeneous plasma by a pulsating point charge $Q \cdot \exp(i\omega t)$, at frequency ω , at a radial distance r from the charge Q is given by:

$$\phi(\omega, r) = \frac{Q}{4\pi\epsilon_0} \frac{2}{\pi} \lim_{Im(\omega) \rightarrow 0} \int_0^{+\infty} \frac{\sin(kr)}{kr} \frac{dk}{\epsilon_l(k, \omega)} \quad (2)$$

where ϵ_l is the longitudinal dielectric function of the plasma, k is the wavelength and ϵ_0 is the vacuum permittivity.

We recall the longitudinal dielectric function ϵ_l for electrostatic waves in an unmagnetized plasma (Krall and Trivelpiece, 1973):

$$\epsilon_l(k, \omega) = 1 + \frac{\omega_{pe}^2}{k^2} \int \frac{k \cdot \nabla_v f_0}{\omega - k \cdot v} dv \quad (3)$$

with f_0 the evdf at equilibrium state, v the electron velocity and ω_{pe} the electron plasma frequency defined by $\omega_{pe} = (n_e e^2 / m_e \epsilon_0)^{1/2}$ where n_e is the electron density, e the electric charge, m_e the electron mass and ϵ_0 the vacuum permittivity.

The longitudinal dielectric function directly depends on the electron velocity distribution function. The evdf typically observed in the solar wind and in magnetospheres can be described as a sum of different evdf as follows (Maksimovic et al., 2005):

$$f_0 = f_{core} + f_{halo} \quad (4)$$

with f_{core} the velocity distribution function of the core electrons, that can be seen as the thermal component, f_{halo} the velocity distribution function of the halo electrons, that can be seen as the suprathermal component. In this work, we have not taken into account other suprathermal electron contributions such as the solar wind strahl (Štverák et al., 2009). While state-of-the-art models of mutual impedance experiments do not enable to model components of the distribution functions that are not symmetric in velocity space (such as the strahl), we later argue and justify that the strahl contribution to the modeling of mutual impedance spectra can be neglected, at least in the limit of the solar wind parameters range close to the perihelion (section 5).

In the literature, f_{core} is usually modeled by a Maxwellian evdf and f_{halo} by a kappa evdf (Lazar et al., 2017). In some cases, f_0 can be directly treated as a single kappa evdf to model both core and halo electrons in a single description on the solar wind for instance (Maksimovic et al., 1997), or in more complex situations as a sum of two kappa evdf as in the magnetosphere of Saturn (Baluku et al., 2011) or in the ionosphere of the comet 67P/Churyumov-Gerasimenko (Clark et al., 2015; Broiles et al., 2016; Myllys et al., private communication).

We use the following notations for a Maxwellian evdf f_{Maxw} and a kappa evdf f_κ :

$$f_{Maxw}(v) = \frac{1}{\pi^{3/2} v_{th}^3} e^{-v^2/2v_{th}^2} \quad (5)$$

$$f_\kappa(v) = (\pi\kappa\theta^2)^{-3/2} \frac{\Gamma(\kappa+1)}{\Gamma(\kappa-1/2)} \left(1 + \frac{v^2}{\kappa\theta^2}\right)^{-(\kappa+1)} \quad (6)$$

where $v_{th} = (k_B T_e / m_e)^{1/2}$ is the electron thermal velocity associated to the electron temperature T_e , k_B the Boltzmann constant, Γ the classical gamma function, $\theta = [(2\kappa - 3)/\kappa]^{1/2} v_{th}$ the generalized thermal speed, with κ is a real number and $\kappa > 3/2$. We remind the reader that the kappa evdf is a generalization of the Maxwellian evdf for $\kappa \rightarrow +\infty$.

In this study, we choose to normalize distances to the Debye length of the Maxwellian evdf $\lambda_{D,Maxw} = (k_B T_e / m_e \omega_{pe}^2)^{1/2}$. As pointed out by Chateau and Meyer-Vernet (1991), the comparison between a Maxwellian and a kappa evdf only makes sense in plasmas characterized by the same density and temperature. In that case, the corresponding Debye length for a kappa evdf is defined as follows:

$$\lambda_{D,\kappa} = \sqrt{\frac{2\kappa - 3}{2\kappa - 1}} \lambda_{D,Maxw} \quad (7)$$

For a collisionless isotropic plasma with a combination of n_M Maxwellian and n_κ kappa evdf, the longitudinal dielectric function ε_l reads (Mace et al., 1999):

$$\varepsilon_l(K, \Omega) = 1 - \sum_{i=1}^{n_M} \frac{Y_i^2}{\Omega_i^2} Z'(Y_i) - \sum_{j=1}^{n_\kappa} \frac{(\kappa_j - 1)^2}{(\kappa_j - 3/2)^2} \frac{Y_j^2}{\Omega_j^2} Z'_{\kappa_j-1} \left[\left(\frac{\kappa_j - 1}{\kappa_j - 3/2} \right)^{1/2} Y_j \right] \quad (8)$$

where:

$$K = k\lambda_{D,ref} \quad (9)$$

$$\Omega = \frac{\omega}{\omega_{pe}} \quad (10)$$

$$\Omega_i = \frac{\omega}{\omega_{pe,i}} \quad (11)$$

$$Y_i = \frac{\Omega_i}{\sqrt{2\mu_i/\tau_i} K} \quad (12)$$

where $\lambda_{D,ref}$ is the Debye length of the hottest electron population. As explained above, if the kappa population is the hottest population, $\lambda_{D,ref}$ is normalized to the corresponding $\lambda_{D,Maxw}$. In addition, we define μ_i (resp. τ_i) the density (resp. temperature) ratio between the hottest population and the i -th population and i.e., $\mu_i = n_{hot}/n_i$ and $\tau_i = T_{hot}/T_i$. Z' and Z'_{κ_j-1} are, respectively, the first derivative of the plasma dispersion function Z (Fried and Conte, 1961) and of the

modified plasma dispersion function Z_κ (Summers and Thorne, 1991). The modified plasma dispersion function Z_κ reads:

$$Z_\kappa(\xi) = \frac{i(\kappa + \frac{1}{2})(\kappa - \frac{1}{2})}{\kappa^{3/2}(\kappa + 1)} {}_2F_1[1, 2\kappa + 2; \kappa + 2; \frac{1}{2}(1 - \xi/i\sqrt{\kappa})] \quad (13)$$

where ${}_2F_1$ is the Gauss hypergeometric function. The main properties of Z and Z_κ can be found in Fried and Conte (1961) and Mace and Hellberg (1995), respectively¹. Using the chosen normalization, Equation (2) that gives the electrostatic potential transmitted in a plasma distance $R = r/\lambda_{D,ref}$ by a pulsating point charge at frequency Ω rewrites (Gilet et al., 2017):

$$\frac{\phi}{\phi_0}(\Omega, R) = \frac{2R}{\pi} \lim_{\text{Im}(\Omega) \rightarrow 0} \int_0^\infty \frac{\sin(KR)}{KR} \frac{1}{\varepsilon_l(K, \Omega)} dK \quad (14)$$

The computation of this radiated electrostatic potential has been carried out using the numerical method described in Gilet et al. (2017) and generalized to a sum of different evdf following Equation (8).

3. ELECTRIC POTENTIAL RADIATED IN A PLASMA WITH SUPRATHERMAL ELECTRONS

In this section, we discuss the radial profile of the electric potential defined in section 2 (Equation 14) for the following electron velocity distribution functions: a kappa evdf (section 3.2) and a sum of a core Maxwellian and a halo kappa evdf (section 3.3). The propagation of the electric potential in the plasma is strongly constrained by the different available linear eigenmodes. We introduce these modes in section 3.1.

3.1. Linear Eigenmodes

We remind the analytic approximation of the linear eigenmodes of the plasma characterized by the different evdf considered in this work (solutions of the dispersion relation $\varepsilon_l(K, \Omega) = 0$) of direct interest in the presence of suprathermal electrons. These modes determine the resonances that shape the mutual impedance spectra. While the longitudinal dielectric function corresponding to a Maxwellian or a kappa evdf has infinite eigenmodes, the least damped modes are the one that contribute most to model the propagation of the electric potential in a plasma. In particular, for a single evdf, the least damped pole, corresponding to Langmuir waves, gives the main contribution to the propagation of the radiated potential in a single electron population plasma, such as a Maxwellian evdf (Chasseriaux et al., 1972; Beghin, 1995). In the large phase velocity limit $\omega/k \gg v_{th}$, with a Maxwellian evdf, the dispersion relation of the Langmuir waves are the following (Krall and Trivelpiece, 1973):

$$\omega_{L,Maxw}(k) = \omega_{pe} \sqrt{1 + 3(k\lambda_D)^2} - i \sqrt{\frac{\pi}{8}} \frac{\omega_{pe}}{(k\lambda_D)^3} e^{-\frac{1}{2(k\lambda_D)^2} - \frac{3}{2}} \quad (15)$$

¹For practical use, we remind that the plasma dispersion function satisfies the differential equation $Z'(y) = -2(1 + yZ(y))$ and derived from the Faddeeva function (or the scaled complex complementary error function): $Z(y) = i\sqrt{\pi} w(y)$

For a single kappa evdf, the Langmuir waves are characterized in the limit $\omega/k \gg \theta$ by the dispersion relation (Mace and Hellberg, 1995):

$$\omega_{L,\kappa}(k) = \omega_{pe} \sqrt{1 + 3(k\lambda_D)^2} - i\pi^{1/2} \frac{\Gamma(\kappa + 1)}{\Gamma(\kappa - 1/2)} \omega_{pe} (2\kappa - 3)^{\kappa-1/2} (k^2 \lambda_{D,\kappa}^2)^{\kappa-1/2} \quad (16)$$

The real frequency (oscillating part) from these dispersion relations are similar, while the damping rate is strongly different and depends on the κ -value. Note that, hereafter, the radiated potentials expressed in a plasma characterized by a Maxwellian or a kappa distribution will be compared for plasmas characterized by the same electron density and temperature. Thus, at equivalent temperature, the Debye length of the kappa evdf is expressed as in Equation (7), so that it is actually smaller than the corresponding Maxwellian Debye length. Note that these analytical expressions are computed within strong approximations (long wavelength limit for instance) that are usually not relevant for the instrumental modeling, as the transmitter-receiver distance can be as small as a few Debye lengths. To go beyond these analytical, though useful, approximations, we also compute numerically the dispersion relations.

Figure 1 (left panel) shows the dispersion relation of the Langmuir pole for a Maxwellian evdf and for different κ -values (from $\kappa = 2$ to 24). From a practical point of view, the position of the Langmuir pole on the real K-space is estimated from the position of the maximum of $Im(1/\epsilon_l(K, \Omega))$, that is plotted in **Figure 1** (right panel). The position of the Langmuir pole projected in the real K-space is similar between the kappa evdf and the Maxwellian evdf, as expected analytically (see Equations 15 and 16). Regarding the damping rate γ , it can be qualitatively constrained by the shape of $Im(1/\epsilon_l(K, \Omega))$ close to the projection of the Langmuir pole on the real K-space. Indeed, the flatter the shape of $Im(1/\epsilon_l(K, \Omega))$, the farther away the pole from the real K-space i.e., the damping rate γ is high.

For a plasma characterized by two different electron populations, such as a sum of two Maxwellian evdf or a mix of a Maxwellian core evdf and a halo kappa evdf, two different modes both strongly contribute to the propagation of the electric potential (Mace et al., 1999; Gilet et al., 2017) namely the (modified) Langmuir mode and the electron acoustic mode. For convenience, we report here only the variation of the real part of the frequency with the wavevector, issued from the

dispersion relations. The damping rate can be found in the hereby mentioned references.

For a plasma modeled by a mix of a Maxwellian core and a halo kappa evdf or by a sum of two Maxwellian evdf (Gilet et al., 2017), in the limit $\omega/k \gg \theta_h \gg v_c$, the dispersion relation of the (modified) Langmuir waves is expressed by:

$$\omega_{L2}(k) = \omega_{pe} \sqrt{1 + 3 \left(\frac{n_h}{n_{tot}} \right)^2 \left(\sqrt{\frac{2\kappa - 3}{2\kappa - 1}} k \lambda_{D,Maxw} \right)^2} \quad (17)$$

In the limit of an intermediate phase velocity i.e., $v_c \ll \omega/k \ll \theta_h$ the dispersion relation of the electron acoustic mode is given by (Mace et al., 1999; Gilet et al., 2017):

$$\omega_{EAW}(k) = \omega_{p,c} \sqrt{1 + 3k^2 \lambda_{D,c}^2 - \frac{1}{\left(\sqrt{\frac{2\kappa - 3}{2\kappa - 1}} k \lambda_{D,Maxw} \right)^2}} \quad (18)$$

We have also computed the useful function $Im(1/\epsilon_l(K, \Omega))$ in a two-electron temperature plasma, in a limit where the electron acoustic and the Langmuir modes co-exist (here $n_h/n_c = 1$, $T_h/T_c = 100$). **Figure 2** shows $Im(1/\epsilon_l(K, \Omega))$ for (i) an evdf modeled by a sum of a Maxwellian core and a kappa halo evdf for different κ -values (κ from 2 to 24) and (ii) an evdf modeled by a sum of two Maxwellian evdf. As expected, $Im(1/\epsilon_l(K, \Omega))$ has two maxima due to the presence of the electron acoustic and the Langmuir modes. For small κ -values, the first pole is not well visible. Indeed, as explained by Mace et al. (1999), for a fixed halo-to-core temperature ratio, the domain of existence of the electron acoustic mode is reduced for lower κ -values.

3.2. Radiated Potential for a Single Kappa Evdf

In order to characterize the effect of suprathermal electrons on the radiated electrostatic potential, we have computed the potential for two frequencies such that no eigenmode propagates in a first case ($\Omega = 0.75$) and a Langmuir mode propagates without being damped much, in a second case ($\Omega = 1.10$). The radial profile of the electrostatic potential, expressed in terms of distance to the transmitter is shown in **Figure 3** for different kappa values ($\kappa = 2, 7$, and 24) and for a Maxwellian evdf ($\kappa \rightarrow \infty$), with equal temperatures (Equation 7). The distances are shown in logarithmic scales.

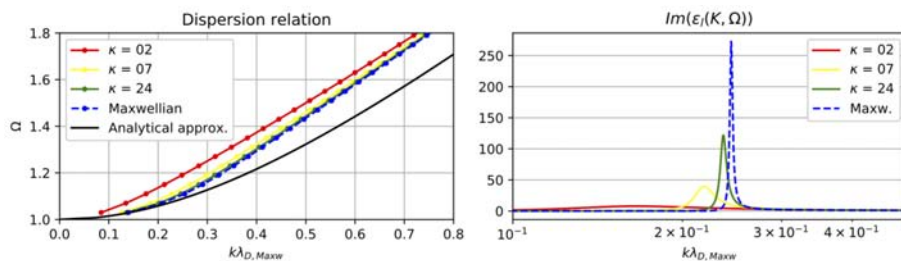


FIGURE 1 | (Left panel) Dispersion relation found by simulation for different κ -values ($\kappa = 2, 7$ and 24) and the Maxwellian evdf ($\kappa \rightarrow +\infty$) with the analytical approximation (black line). **(Right panel)** $Im(1/\epsilon_l(K, \Omega))$ for the same evdf.

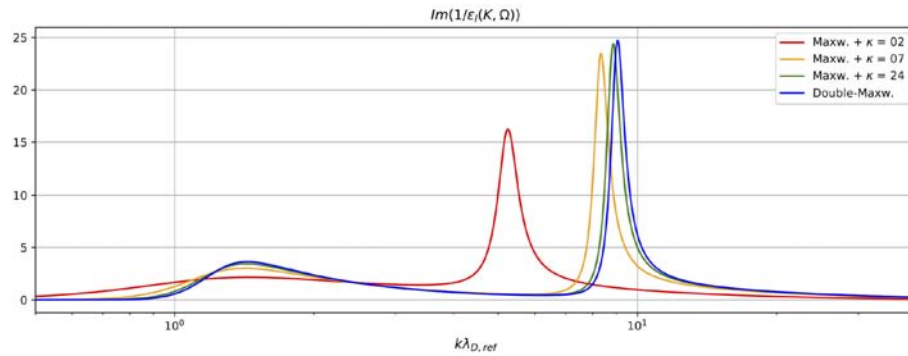


FIGURE 2 | $\text{Im}(1/\epsilon_l(K, \Omega))$ for $\Omega = 1.10$ in a plasma modeled by a mix of a core Maxwellian evdf and a halo kappa evdf for several κ -values ($\kappa = 2, 7, 24$) and a double Maxwellian evdf, where the modified Langmuir mode and the electron acoustic mode co-exists (here: $n_h/n_c = 1$ and $T_h/T_c = 100$). $\lambda_{D,Ref}$ is the Debye length of the halo Maxwellian evdf.

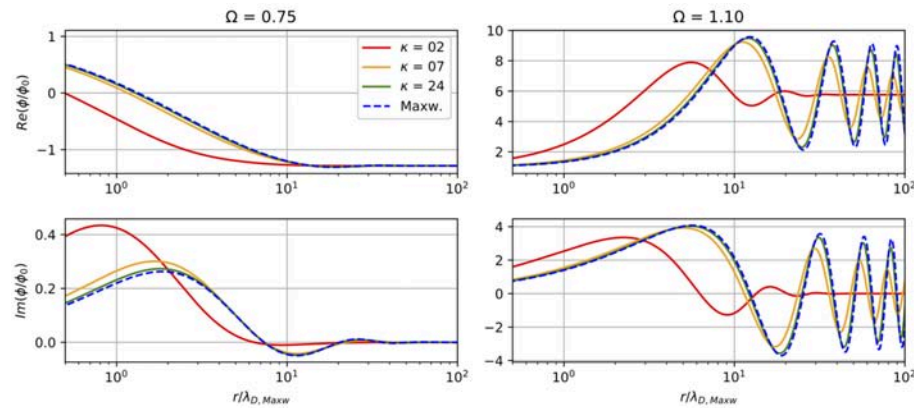


FIGURE 3 | Radiated electrostatic potential for $\Omega = 0.75$ (left column) and $\Omega = 1.10$ (right column) compared to the distance to the transmitter normalized by $\lambda_{D,Maxw}$ in logarithmic scale: the real part (first panel) and the imaginary part (second panel) for different kappa evdf (here $\kappa = 2, 7$ and 24) in colored continuous lines and a Maxwellian evdf ($\kappa \rightarrow +\infty$) in dashed line.

Note that for the two frequencies Ω , the real part of ϕ/ϕ_0 tends to the inverse of the cold plasma dielectric constant $\epsilon_c = \frac{1}{1-\Omega^2}$ (here $\epsilon_c^{-1} = -1.29$ for $\Omega = 0.75$ and $\epsilon_c^{-1} = 5.76$ for $\Omega = 1.10$) and the imaginary part tends to 0, as expected (Beghin, 1995; Gilet et al., 2017).

At frequencies higher than the electron plasma frequency, here $\Omega = 1.10$ (right column), the real and the imaginary part of the radiated potential oscillate. The radiated potential in a plasma modeled by a kappa evdf tends to the potential of the Maxwellian evdf when κ increases (here $\kappa > 10$) as expected. However, for the low κ -values, the radiated potential is more damped. This is explained by the higher damping rate γ (section 3.1) for the evdf characterized by the presence of suprathermal electrons. Moreover, the wavelength of the oscillations decreases (from $\sim 25R$ to $\sim 14R$) while the suprathermal electrons contribution increases, as expected from the linear theory of Langmuir waves in a kappa distribution plasma².

²Note that the wavelength computed from the analytical dispersion relation ($\lambda = 2\pi/K$) is close to the wavelength computed numerically and corresponding to the oscillations of the modeled radiated potential, as expected.

For frequencies lower than the electron plasma frequency ($\Omega = 0.75$ in **Figure 3**, left panel), the radiated potential does not oscillate because no eigenmode exists at this frequency range. This has a strong implication on the mutual impedance spectrum in particular when the transmitter-receiver distance is short compared to the Debye length that is developed in section 4.

3.3. Radiated Potential for a Mix of Kappa and Maxwellian evdf

We have also investigated the radial variation of the radiated electric potential injected in a plasma modeled with a mix of a core Maxwellian evdf and a halo kappa evdf, as typically observed in the solar wind plasma. In this section, all distances are normalized to the Debye length of the Maxwellian evdf corresponding to the Debye length of the kappa evdf (see Equation 7). The computed potential is illustrated in **Figure 4** in a region where the electron acoustic mode exists (here $n_h/n_c = 0.4$ and $T_h/T_c = 100$) for different κ -values.

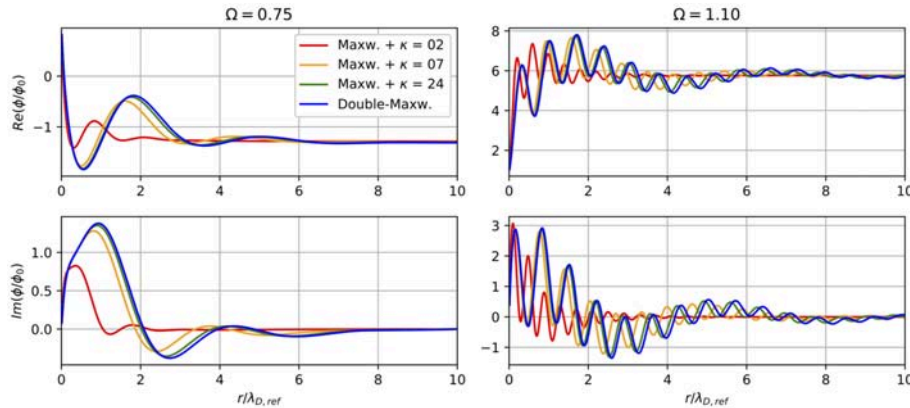


FIGURE 4 | Radiated electrostatic potential compared to the distance to the transmitter normalized by $\lambda_{D,ref}$ for $\Omega = 0.75$ (left column) and $\Omega = 1.10$ (right column) with the real part (first line) and the imaginary part (second line) in a two-electron temperature plasma ($\mu = 0.4$ and $\tau = 100$), for (i) a mix of a core Maxwellian and a halo kappa evdf (here $\kappa = 2, 7, 24$) and (ii) a sum of two Maxwellian evdf (blue line).

First, at frequencies higher than the electron plasma frequency (here $\Omega = 1.10$), the radiated potential is characterized by a superposition of two characteristic waves due to the transmission of both electron and Langmuir fluctuations (section 3.1), as been observed in **Figure 4** (right column). In this case, the Langmuir wavelength is larger than the electron acoustic wavelength. For both oscillations, the waves are more damped when there are more suprathermal electrons in the plasma (i.e., for decreasing κ), as expected from a large Landau damping at small κ .

Second, at frequencies smaller than the electron plasma frequency (here $\Omega = 0.75$), contrary to the potential radiated in a plasma with a single evdf, the potential oscillates due to the electron acoustic mode. The potential is more damped when the suprathermal part increases (i.e., κ -value decreases). Note that this oscillation is strongly damped, though, so that we do not expect the signal propagating further in the plasma. This means that in the case of a receiver located far (in terms of ion acoustic wavelengths) from the transmitter, we do not expect a strong signature in the mutual impedance spectra, while the instrument shall be sensitive to the ion acoustic mode adapted to the transmitter-receiver distance, i.e., we expect the mutual impedance spectra to exhibit the signature of the ion acoustic mode which wavelength is twice the transmitter-receiver distance.

3.4. Mutual Impedance Responses

The potential modeled in the previous section is used to compute the mutual impedance response. Indeed, the transmitters inject an oscillating current $I(\Omega)$ at a given frequency while the receivers measure the (complex) amplitude of the electric potential $V(\Omega)$ at the same frequency. The mutual impedance $Z(\Omega) = \Delta V(\Omega)/I(\Omega)$ is then directly related to the difference between the electric potential $\Delta V(\Omega) = V_{R_2}(\Omega) - V_{R_1}(\Omega)$, radiated by the different emitters at frequency Ω and measured by two receivers R_1 and R_2 . To isolate the effect of the plasma to the potential radiated by the emission part of a mutual impedance probe, we work with the mutual impedance spectrum normalized

to the spectrum that is obtained in vacuum

$$H(\Omega) = \frac{Z}{Z_0} = \frac{V_{R_2}(\Omega) - V_{R_1}(\Omega)}{V_{R_2,0} - V_{R_1,0}} \quad (19)$$

where Z and Z_0 represent the mutual impedance of a probe surrounded by a plasma and by the vacuum, respectively, and V_{R_i} (resp. $V_{R_i,0}$) is the voltage measured by the receiver R_i in the plasma (resp. in vacuum):

$$V_{R_i}(\Omega) = \frac{1}{4\pi\epsilon_0} \sum_{j=1} \frac{\phi}{\phi_0}(\Omega, d_{ij}/\lambda_{D,ref}) \frac{q_j}{d_{ij}} \quad (20)$$

$$V_{R_i,0} = \frac{1}{4\pi\epsilon_0} \sum_{j=1} \frac{q_j}{d_{ij}} \quad (21)$$

where q_j is the charge of the j th transmitter and d_{ij} is the distance between the receiver R_i and the j th transmitter, ϕ and ϕ_0 are the electric potential radiated by a pulsating point charge embedded in the plasma or within vacuum, respectively.

The electron plasma frequency is located in the close vicinity of the maximum amplitude of the mutual impedance response (Storey et al., 1969; Chasseriaux et al., 1972). The total electron density, n_{tot} , is then determined from the electron plasma frequency f_{pe} with $n_{tot} = (f_{pe}/8.98)^2$ (n_{tot} is expressed in cm^{-3} and $f_{pe} = \omega_{pe}/2\pi$ in kHz).

4. APPLICATION TO THE BEPICOLOMBO MUTUAL IMPEDANCE PROBE AM²P

In this section, we apply the modeling of the electric potential radiated in a plasma with suprathermal electrons, described previously, to the computation of synthetic mutual impedance spectra. We aim at characterizing the effect of suprathermal electrons on instrumental response of the mutual impedance probe AM²P of the Plasma Wave Investigation (PWI) consortium (Kasaba et al., 2010) onboard the Mercury

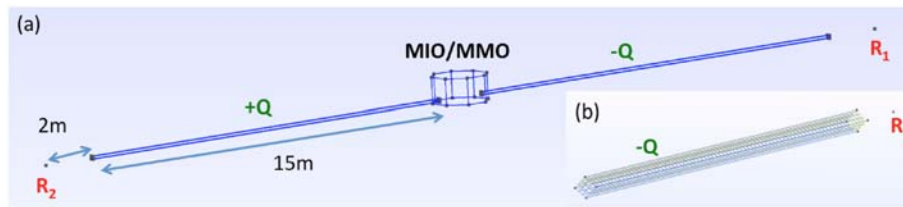


FIGURE 5 | (a) AM²P geometry consists on two 15 m-antennas on both sides of the MMO spacecraft with two receivers, R_1 and R_2 , located at 2 m of each end of booms, **(b)** example of the meshing of one AM²P antenna.

Magnetospheric Orbiter (MIO/MMO) of the BepiColombo mission. The PWI/AM²P experiment will measure the plasma bulk properties of the Mercury magnetospheric and Solar wind plasma such as the electron density (in the 0.02 to 180 cm^{-3} range, corresponding to f_{pe} from 0.7 to 120 kHz) and the electron temperature, in a range which depends on plasma conditions (Trotignon et al., 2006). The BepiColombo spacecraft has been launched successfully in October 2018, for an interplanetary cruise phase of 7.2 years (until December 2025) with one Earth flyby, two Venus and six Mercury flybys before the nominal mission science operations performed for one and a half Earth year (about 6 Hermean years) and a planned extension of one Earth year, corresponding to 4 extra Hermean years).

The MIO/MMO spacecraft will have an elliptic polar orbit of $400 \times 11,824 \text{ km}$ (Benkhoff et al., 2010). From the observations of the MESSENGER mission (Johnson et al., 2012), the interplanetary and Hermean magnetic fields are such that the electron cyclotron frequency is expected to be negligible compared to the electron plasma frequency for low latitudes or high enough distances from Mercury. The modeling of the electric radiated potential described in this paper is only valid in an unmagnetized plasma, i.e., where the electron-cyclotron frequency f_{ce} is negligible compared to the electron plasma frequency f_{pe} , therefore, we hereafter focus on the AM²P modeling in the solar wind plasma and in the Hermean magnetosphere far from the cusps. Other analysis methods shall be considered (Béghin et al., 2017) or developed in strongly magnetized regions.

4.1. PWI/AM²P Antenna Configuration

The PWI/AM²P quadrupolar probe consists of (i) two transmitting 15m-antennas of 1cm-diameter located on both sides of the MIO/MMO spacecraft and (ii) two receivers located at 2m of the end of the transmitting antennas (Figure 5). In this model, the transmitting antennas have been discretized in about thousand rectangular sub-elements, with the center of each sub-element considered as a pulsating point charge, while the receiving antennas are considered as being punctual. This experiments works in the so called Double-Wire (or push-pull) mode for which the pulsating charge on one transmitting antenna is opposite to that of the second transmitting antenna, in other words they are in phase opposition. Given this geometry and charge configuration, the expected mutual impedance is

modeled using Equation (19), combined with Equations (20) and (21).

4.2. Modeling of the AM²P Mutual Impedance Spectra

In the following, the synthetic instrumental response of PWI/AM²P is computed for (i) a single kappa evdf (section 4.2.1) and (ii) a mix of a halo kappa evdf and a core Maxwellian evdf (section 4.2.2). In this section, we consider a large range of plasma parameters in order to characterize the effect of suprathermal electrons in different regimes. We will focus on the plasma conditions in the solar wind and at Mercury expected to be encountered by the MIO/MMO spacecraft in section 5.

4.2.1. AM²P Spectra With a Single Kappa evdf

We have modeled the PWI/AM²P mutual impedance response for different κ -values, as well as for a Maxwellian evdf ($\kappa \rightarrow +\infty$) for direct comparison and validation. The mutual impedance response is computed for different plasma conditions characterized by the (equivalent Maxwellian) Debye length of the hottest electron population (from 30 cm to 5 m, renormalized by the corresponding Maxwellian evdf). The results are shown in Figure 6.

First, in the limit of the Debye length is much smaller than the transmitter-receiver distance, the mutual impedance spectra are similar whatever the presence of suprathermal electrons (top left and right panels, corresponding to $\lambda_{D,Maxw} = 30 \text{ cm}$ and $\lambda_{D,Maxw} = 1 \text{ m}$). In this regime, the mutual impedance measurement principle is therefore transparent to the presence and nature of suprathermal electrons and robust in determining the total electron plasma density.

Second, when the Debye length is slightly smaller, the mutual impedance spectra is flatter for high κ -values or a Maxwellian evdf than for low κ -values (bottom left and right panels, corresponding to $\lambda_{D,Maxw} = 2 \text{ m}$ and $\lambda_{D,Maxw} = 5 \text{ m}$). Moreover, the maximum of the amplitude is shifted compared to the total electron plasma frequency for high κ -values or a Maxwellian evdf. In this regime, the presence of suprathermal electrons enables to detect the total plasma frequency on the mutual impedance spectra. This counter-intuitive result must be balanced by the fact that the mutual impedance spectra is computed for a smaller Debye length when the κ -value decreases. The comparison needs to be performed in the same plasma i.e., same electron density and electron temperature. Note that the

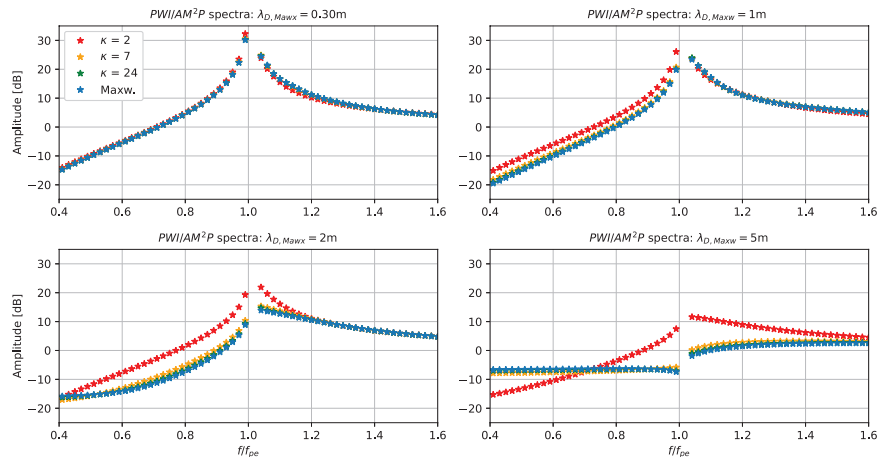


FIGURE 6 | Modeled PWI/AM²P mutual impedance spectra for different $\lambda_{D,Maxw}$ (from 30 cm to 5 m) in a plasma modeled by a kappa evdf ($\kappa = 2, 7, 24$) and a Maxwellian evdf ($\kappa \rightarrow +\infty$). The amplitude is expressed in $20\log_{10}$ and the frequency is normalized by the electron plasma frequency.

shape of mutual impedance response of a Maxwellian for $\lambda_{D,Maxw} = 2$ m is similar to the response for a kappa evdf for $\lambda_{D,Maxw} = 5$ m. Therefore, it is not possible to characterize the suprathermal electrons from the AM²P spectra.

4.2.2. AM²P Spectra With a Mix of a Halo Kappa and a Core Maxwellian evdf

To go beyond, we consider a plasma with a mix of a halo kappa and a core Maxwellian evdf, as observed in the solar wind by Pierrard et al. (2016). The AM²P spectra have been computed in a large range of plasma parameters: the core-to-total density ratio n_c/n_{tot} varies from 0.1 to 0.9 and the halo-to-core temperature ratio T_h/T_c varies from 10 to 500 with the same Debye length $\lambda_{D,Maxw} = 4$ m. This is reported in **Figure 7**, where the halo-to-core temperature ratio increases from left to right, while the density of the core electrons increases from bottom to top.

First of all, when the density of the core population is much higher than one of the halo (top panels), the mutual impedance spectrum is close to what is observed in a plasma modeled by a single evdf (**Figure 3**, top panels). Only one resonance appears close to the *total* electron frequency. In this limit, the response is independent to the κ -values of the halo evdf.

Second, when the plasma contains as many core electrons as halo electrons (middle row) or when the electron density is dominated by the halo part (third row), the shape of the mutual impedance spectra depends on the κ -value. As seen in the previous section, the resonance at the total plasma frequency is flatter when the kappa evdf tends to the Maxwellian evdf. With the Debye length considered here, the total electron density can be estimated for all κ -values. When the halo-to-core temperature ratio increases, a second resonance appears close to the core plasma frequency (blue vertical dotted line). This resonance is more pronounced when the halo-to-core temperature ratio increases whatever the κ -value. At a given halo-to-core temperature the amplitude of the electron acoustic mode increases with κ -value. This could be explained by the

decay of the electron-acoustic mode domain of existence, with a fixed halo-to-core temperature ratio, in a presence of a mix of kappa and a Maxwellian evdf when the suprathermal electron part increases (Mace et al., 1999).

5. DISCUSSION

In the previous section, the mutual impedance spectra was modeled in a large domain of plasma parameters to characterize the effect of suprathermal electrons in the mutual impedance measurement. In this section, the AM²P spectra is computed in the plasma conditions expected to be encountered by the MIO/MMO spacecraft: in the Hermean magnetospheric plasma (section 5.1) and in solar wind plasma close to the perihelion and the aphelion of Mercury (section 5.2). For that, we used the evdf found by fitting method with *in-situ* evdf measurement of the solar wind plasma (Maksimovic et al., 1997; Pierrard et al., 2016) and the energetic particle measurement of Mariner 10 during a flyby in the Hermean magnetospheric plasma (Baker et al., 1986).

5.1. AM²P Spectra in the Hermean Magnetosphere

A large part of the elliptical orbit of MIO/MMO will be in the Hermean magnetosphere. In order to characterize the effect of the magnetospheric plasma in the AM²P spectra, we used the observations of the electron density and the electron temperature measured by Mariner 10 during a flyby in the Mercury magnetospheric plasma in the nightside of Mercury with the closest approach at 700 km of the surface. These measurements are summarized in Baker et al. (1986). The modeling of the mutual impedance response in the electrostatic limit is valid due to the fact that the cyclotron frequency was negligible (around 3 kHz) compared to the total plasma frequency (around 20 kHz). Different electron populations should be observed in the Hermean magnetosphere especially: (i) an electron population from the solar wind origin (n_{SW} from 7 to 12 cm⁻³, T_{SW} from 22

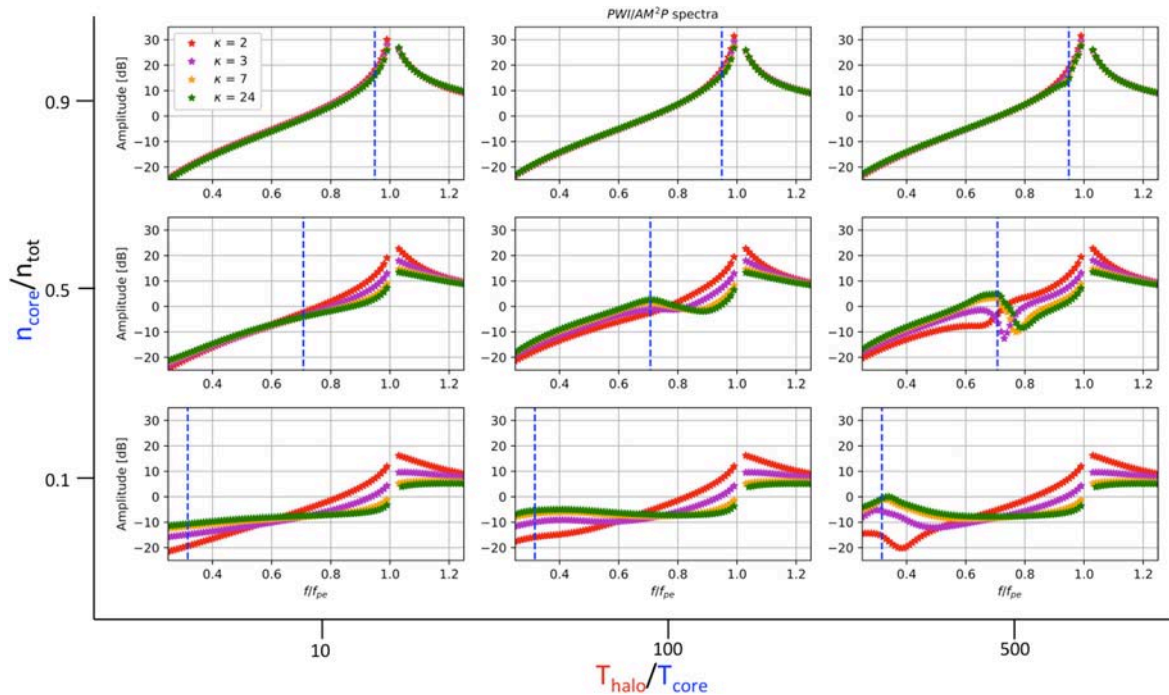


FIGURE 7 | Modeled PWI/AM²P mutual impedance spectra for $\lambda_{D,ref} = 4$ m in a plasma modeled by a mix of a core Maxwellian evdf and a halo kappa evdf ($\kappa = 2, 3, 7, 24$). The amplitude is expressed in $20\log_{10}$ and the frequency is normalized by the electron plasma frequency. The core plasma frequency is shown by the blue vertical dotted lines.

to 40 eV) and (ii) an electron population from the magnetosheath (n_{MAG} from 3 to 7 cm^{-3} , T_{MAG} from 12 to 40 eV). A mixed of the two populations can be observed in the magnetospheric plasma. We choose to model the two electron populations both by a Maxwellian evdf. The Debye length of the solar wind electron population is characterized by (resp. magnetosheath) 10.0 m $< \lambda_{D,SW} < 17.7$ m (resp. 9.7 m $< \lambda_{D,MAG} < 27.0$ m). Several examples of the modeled AM²P spectra with different configuration of the mix of the two electron populations are shown in **Figure 8**. For the considering cases, the resonance above the total electron plasma frequency is particularly flat, due to the large Debye length of the two electron populations compared to the transmitter-receiver distance. Therefore, the detection of the total electron plasma frequency, and therefore the measurement of the electron density will be challenging in the Hermean magnetosphere. Moreover, the resonance close to the plasma frequency corresponding to the electrons from the magnetosheath is not visible due to the fact that the temperature ratio is too low (see **Figure 7**). Therefore, the presence of two electron populations might not be observed by AM²P in this regime of parameters.

5.2. AM²P Spectra in the Solar Wind Plasma Close to the Perihelion and Aphelion of Mercury

Mercury has the largest planetary orbital eccentricity in the Solar system. The distance to the Sun varies from 0.31 AU at perihelion

to 0.47 AU at aphelion. We modeled the AM²P spectra in the solar wind plasma for both heliocentric distance.

First, close to the perihelion at 0.35 AU, the evdf of the solar wind has been characterized by a mix of a halo kappa evdf and a Maxwellian core evdf with a halo-to-core density ratio n_{halo}/n_{core} equals to 0.03 ($n_{core}/n_{tot} = 0.97$) and a halo-to-core temperature ratio T_{halo}/T_{core} equals to 3.36 with $\kappa_h = 7.54$ (Pierrard et al., 2016). The Debye length of the core population (resp. halo) is $\lambda_{D,core} = 3.73$ m (resp. $\lambda_{\kappa,halo} = 33.4$ m). The corresponding mutual impedance spectra in the solar wind at 0.35 AU is shown in **Figure 9** (red line, right panel). In order to characterize the effect of the suprathermal electrons in the solar wind in the AM²P spectra, the AM²P spectra has been modeled with only the core Maxwellian evdf (blue dotted line, right panel). We observed that the AM²P spectra modeled by a sum of a core Maxwellian and a halo kappa evdf (red curve) and only with the core Maxwellian (blue curve) are similar. Also, the AM²P spectra is flat (large resonance spectral signature) with a spectral peak ($\Omega = 1.25$) shifted with respect to the plasma frequency ($\Omega = 1$), while the cut-off frequency enables to retrieve efficiently the plasma frequency. Therefore, close to the perihelion of Mercury, the AM²P experiment is robust to the presence of suprathermal electrons, seen as the halo part of the evdf, in the solar wind when determining the total electron density.

Second, the AM²P spectra has been modeled in the solar wind plasma at 0.5 AU, close to the aphelion of Mercury (0.47 AU). The halo-to-core density ratio n_{halo}/n_{core} is equals to 0.04, the halo-to-core temperature T_{halo}/T_{core} is equals to 4.10 and $\kappa_h =$

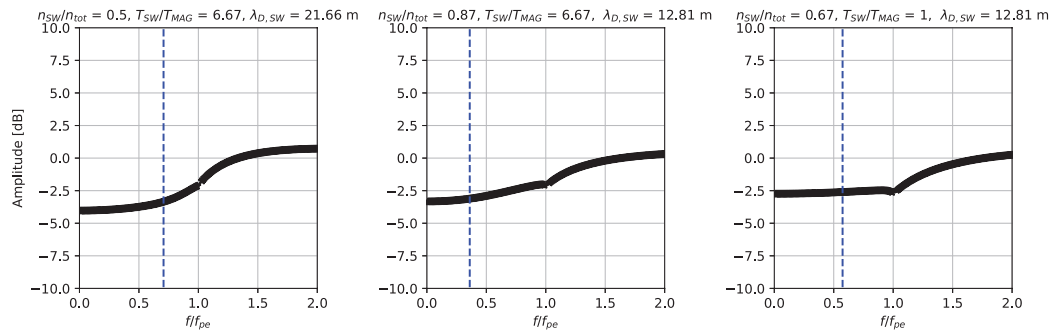


FIGURE 8 | Modeled AM²P spectra in three different plasma with a mix of electrons from the solar wind origin and the electrons from the magnetosheath, both modeled by a Maxwellian evdf. The amplitude is expressed in $20\log_{10}$ and the frequency is normalized by the electron plasma frequency. The core plasma frequency is shown by the blue vertical dotted lines.

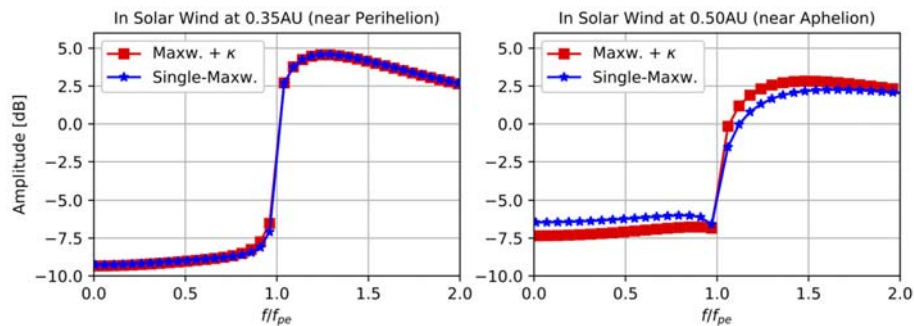


FIGURE 9 | Modeled PWI/AM²P mutual impedance spectra for (i) a core Maxwellian evdf and a halo kappa evdf (red squared line) and (ii) only the core Maxwellian evdf (blue asterisk line) in the *in-situ* measured solar wind plasma at 0.35 AU (near Mercury perihelion, left panel) and 0.5 AU (near Mercury aphelion, right panel). The amplitude is expressed in $20\log_{10}$ and the frequency is normalized by the total electron plasma frequency.

6.89 with $\lambda_{D,halo} = 46.37$ m, $\lambda_{D,core} = 5.24$ m (Pierrard et al., 2016). The modeled spectra is shown in **Figure 9** (right panel). Compared to the AM²P spectra close to the perihelion (left panel), the AM²P spectra at 0.5 AU is flatter. The maximum of amplitude is around 3 dB. This maximum is located far from the plasma frequency ($\Omega = 1.5$) but the plasma frequency can be retrieved by the cut-off frequency. Due to the instrumental noise, we expect that the signal shall be measurable, with a low signal-to-noise ratio. Contrarily to the plasma conditions in the solar wind near perihelion, the shape of the spectra is affected by the suprathermal electrons modeled by a halo kappa evdf. Indeed, the spectra corresponding to the single-Maxwellian core (blue asterisk line) is different to the spectra modeled by the mix of the core Maxwellian and the halo kappa. However, the shape is closely similar which do not enable to separate the two electron populations and therefore it do not provided a measurement of the suprathermal electrons.

In this study, we have assumed that the solar wind strahl can be neglected. At the location of the perihelion (0.31 AU) and the aphelion (0.45 AU) of Mercury, the strahl contribution represents around 2–3% of the total electron density, less than the halo that represents from 8 to 10% (Maksimovic et al., 2005). Also the

“equivalent” strahl temperature would be higher than the halo one. Since in similar conditions expected to be encountered by BepiColombo, the halo contribution to the mutual impedance spectra is found to be negligible close to the perihelion, therefore, we expect the strahl contribution to the mutual impedance spectra to be even less significant than the halo part itself. However, the halo evdf can modify the shape of the spectra close to the aphelion. Therefore, if the strahl might slightly and marginally affect the mutual impedance spectra, we expect it to be within 1 dB which is hardly detectable experimentally.

6. CONCLUSION

Mutual impedance experiments strongly depend of the electron velocity distribution function (evdf) encountered in the *in-situ* observed space plasma. This study illustrates the influence of suprathermal electrons on the instrumental response of mutual impedance experiments in the interplanetary plasma where the Debye length is of the order of the transmitter-receiver distance. Suprathermal electrons are observed in the Solar system as in the solar wind (Maksimovic et al., 1997), in the ionosphere of the comet 67P/Churyumov-Gerasimenko (Clark et al., 2015; Broiles

et al., 2016; Myllys et al., private communication) and in the Hermean magnetosphere (Christon, 1987). These electrons are usually modeled by (i) a kappa (Maksimovic et al., 1997) or a mix of core Maxwellian evdf and a halo kappa evdf (Pierrard et al., 2016). Thus, we have modeled the longitudinal dielectric function and the electrostatic radiated potential in a plasma modeled by these two different evdf, using and extending the numerical method developed in Gilet et al. (2017). We apply the modeling in the case of the mutual impedance experiments PWI/AM²P onboard the MIO/MMO spacecraft of the BepiColombo mission, successfully launched in October 2018. First, we show that for a single evdf such as a kappa evdf, the radiated potential is more damped and the wavelength is smaller with the presence of suprathermal electrons. For the same electron temperature, the (Langmuir) resonance close to the electron plasma frequency is more visible on mutual impedance spectra than for a Maxwellian evdf when the Debye length increases and is of the order of the transmitter-receiver distance. When the plasma is modeled by a core Maxwellian evdf and a halo kappa evdf, as for a sum of two Maxwellian evdf (Gilet et al., 2017), an other resonance appears before the total plasma frequency due to the existence of the electron acoustic mode in a certain domain of the core-to-halo density and temperature. When the halo evdf is modeled by a kappa evdf with a low κ -value, the resonance due to the electron acoustic mode is less visible on the mutual impedance spectra. Second, we apply the modeling in a more realistic plasma in order to characterize the robustness of the experiment in the Hermean magnetospheric and the solar wind plasma. We show that the halo component of the evdf typically observed in the solar wind is neglected by the AM²P experiment close the perihelion but it can slightly affect the spectra close to the aphelion. Moreover, the AM²P experiment operating in Double-Wire (push-pull) mode should be in the limit of the measurement of the electron density when operating in the low density Hermean magnetospheric plasma. We expect the mutual impedance spectra acquired in these regions to be rather flat so that the expected resonance close to the plasma frequency might not be clearly visible in the low-density plasma surrounding Mercury. Therefore, the detection of the plasma frequency might be challenging for the AM²P experiment in such regions. Measurements of the plasma bulk properties from SORBET and the Langmuir Probes might cover the range of the electron density measurements in these regions. Note that the quasi-thermal noise spectroscopy is also

sensitive to the presence of suprathermal electrons modeled by a kappa evdf (Le Chat et al., 2009). In the contrary, this study shows that mutual impedance spectra acquired in the solar wind close to Mercury where MIO/MMO shall spend most of the operating time (either the free solar wind, the magnetosheath, or the mixing layer between the solar wind and the Hermean plasmas) will give access to the plasma density. In particular, the modeling of the AM²P mutual impedance spectra described in this paper shows that, in the solar wind plasma, the mutual impedance cut-off frequency will represent a fast and efficient estimation of the plasma frequency, and therefore of the plasma density, which represent direct useful practical input for the future data processing of the AM²P instrument. This work should be used also in the future, for the mutual impedance experiment RPWI/MIME onboard the JUICE spacecraft. This experiment will operate in the Jovian system in order to constrain the plasma bulk properties in the Jupiter magnetosphere and in particular in the ionosphere of Ganymede.

AUTHOR CONTRIBUTIONS

NG: performed the analysis and wrote the paper; PH: Ph.D. supervisor, Lead Co-Investigator of AM²P, supervised the paper; GW: helped about the modeling of the mutual impedance response; MM: helped about the suprathermal electrons; OR: helped about the plasma physics; CB: helped about the background of the story of the mutual impedance experiment; J-LR: Principal Investigator of the RPWI/MIME experiment onboard JUICE, helped about the principle of the experiment.

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Experiments With Plasmas Artificially Injected Into Near-Earth Space

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Plasma injection experiments in space are being ordered according to five aspects: (1) Diagnostics of electric fields, (2) Coupling to the ionosphere, (3) Interactions with the solar wind, (4) Modification experiments, and (5) Special physical processes. Historically first were releases of neutral gases with the aim to measure atmospheric parameters. They were soon followed by plasma injections applied to the measurement of plasma flows and parallel electric fields. Long-range coupling to the environment was a most important aspect of the plasma releases. It concerned, on the one hand, the need for corrections of the derived diagnostic parameters and, on the other hand, the understanding of the formation of the ubiquitous striations and deformations of the plasma clouds. A special application was the investigation of cometary interactions by releases in the solar wind. Modification experiments in the ionosphere were done intentionally or occurred as byproducts of rocket launches or other activities. A particular goal was to trigger natural large-scale ionospheric instabilities like equatorial spread F in order to improve the understanding of the natural phenomena. Large-scale plasma injections in the magnetosphere have been performed in order to change the conditions of wave-particle interactions and potentially trigger observable effects. Special goals were so-called skidding experiments and testing Alfvén's critical ionization velocity effect. In this review, we will emphasize the principle objectives and illustrate the results from selected experiments.

Keywords: barium clouds, coupling to ionosphere, modification experiments, artificial comets, critical ionization velocity, auroral acceleration, auroral stimulations

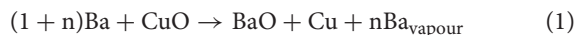
DIAGNOSTICS OF ELECTRIC FIELDS BY TRACING VISIBLE PLASMA CLOUDS

It was Bates (1950), who made the first proposal to release metallic and molecular vapor clouds in the upper atmosphere as a tool to measure atmospheric parameters and excitation processes. Beginning in the mid-fifties the idea was taken up by various groups in the United States (Edwards et al., 1956), Australia (Groves, 1960; Rees, 1961), and France (Blamont et al., 1960). A large number of elements were employed in chemical releases, metallic atoms such as Li, K, Na, Sr, Ba, Al, and compounds such as NO, NO₂, SF₆, AlO, and BaO. The main goal was the investigation of diffusion processes, wind profiles, density, turbulence, and temperature. The technique had the advantage of relative technical simplicity, short preparation times, and many novel insights into a hitherto little known territory. A short overview can be found in Harang (1969).

This review will not deal with the accomplishments pertaining to the neutral atmosphere, but concentrate on the exploration of the near-Earth plasma environment. It soon turned out that the injection barium vapor clouds was the most efficient tool. The first successful releases of barium were made by Armstrong (1963) at the Weapons Research Establishment in Woomera, Australia. The releases were carried out at heights of only 100 km and led to barely detectable ion emissions not yet exhibiting strong coupling to the magnetic and electric fields. However, it was to be expected that the groups working with chemical releases from sounding rockets would soon widen the application to the plasma realm. The main push came from a completely different direction.

It was his study of cometary plasma tails as natural probes of the state of the interplanetary medium what let Ludwig Biermann at the Max Planck Institute for Physics and Astrophysics in Munich propose to perform release experiments suited to create conditions similar to those found in cometary ion tails (Biermann et al., 1961). His proposal fortuitously coincided with the decision of the Federal Republic of Germany in 1961 to enter space research and to charge Biermann's Institute for Astrophysics with that task. Reimar Lüst, although a theorist, was asked to form the first German team dedicated to space research. Taking up Biermann's proposal of creating an "artificial comet" as the first goal, promised to be a quick start. It was of course not regarded necessary to employ materials as found in comets, such as CO^+ , N_2^+ , etc. Materials with low photo-ionization potential and resonance lines in the visible spectral range, such as Sr and Ba, appeared to be the best choice. It was of invaluable help that Jacques Blamont offered the group of Reimar Lüst space on his sounding rockets. The first experiments in 1963 with mixtures of barium oxide and aluminum turned out not to produce significant amounts of free Ba atoms (Föppl et al., 1965). In late 1964, however, at experiments carried out on flights from Hammaguir in the Sahara desert, the observers saw for the first time substantial emissions of Ba I and Ba II (Föppl et al., 1967).

The first experiments were typically performed at heights of 150 km and primarily dedicated to finding the most efficient chemical reaction. Instead of the conventional thermite mixtures, using the reaction:



with n chosen above the stoichiometric equivalent to the amount of CuO, turned out to be the best choice. It became the standard mix for practically all subsequent experiments. A second most helpful result from these first experiments was the finding that photoionization of atomic barium was much faster than calculated in Föppl et al. (1965). Later experiments at higher altitude allowed to determine the true time-scale of nearly 30 s. The explanation was found by Haser (1967), namely a two-step process involving excitation of two metastable levels and ionization from there by the more intense solar UV emissions (Drappatz, 1972; Carlsten, 1975).

At the same time, Haerendel et al. (1967) addressed the relation between the observed motion of a Ba^+ cloud and the transverse electric field. The essential point is that the

modification of the electric field by the locally enhanced Pedersen conductivity is reduced by current exchange between the more highly conducting lower ionosphere and the weakly conducting cloud. Crucial is the ratio, $\lambda^* = \Sigma_{P2}/\Sigma_{P1}$, between the integrated Pedersen conductivity, Σ_P , inside the whole flux tube pervading the cloud (index 2) and that at the outside (index 1). The ratio κ_i between the gyro- and collision frequencies for the barium ions is typically well above unity for clouds generated in the middle or even upper F region. For the E region $\kappa_i \approx 1$. **Figure 1** depicts the situation. It shows a deviation of part of the E region Pedersen currents through the barium cloud by exchanging field-aligned currents. Thereby the contribution of the E layer to the electric polarization field is being reduced mitigating its enhancement by the presence of the barium plasma. Since the cloud is finite in three dimensions, also the Hall conductivity is of importance. However, in most cases it can be safely assumed that the Ba-cloud adds very little to it. Furthermore, the neutral wind velocity at cloud level, v_n , enters into the relation derived for the transverse electric field, $E_\perp$, the quantity of prime interest. Assuming that, owing to the high parallel conductivity, magnetic field lines are equipotentials, is a good approximation within the ionosphere. The relation connecting the observed cloud motion, $V_\perp$, with the transverse electric field:

$$E_\perp^0 \approx \frac{1 + \lambda^*}{2} \frac{B}{c} \left[e_B \times V_\perp + \frac{1}{\kappa_i} (V_\perp - v_{n\perp}) + \frac{\lambda^* - 1}{\lambda^* + 1} v_{n\perp} \times e_B \right] \quad (2)$$

could be used in two ways. On the one hand, it allowed for the first a reliable determination of magnetospheric electric fields and, on the other hand, served for estimating the contribution of polarization fields. The latter are the cause of the frequently evolving distortions of the clouds by creating divergences of the Pedersen current in the E region and secondary conductivity changes thus modulating the $E \times B$ drift of the barium ions. Of course, the external electric field can be quite inhomogeneous from the outset, in particular in the auroral magnetosphere. In the following we will first concentrate on the first products of the barium cloud technique, the insights obtained into the magnetospheric electric fields, and address distortions and fine structures in section Distortions and Striations.

In the late 60s and early 70s direct electric field measurements were still in the development phase. Electric fields derived from Ba cloud motions provided the first trustworthy information about the convection of the magnetospheric plasma, albeit restricted to the times of sunrise and sunset due to the experimental conditions. A summary of the first years' experiments, not only by the Max Planck group, but augmented by similar experiments taken up by the University of Alaska (Wescott et al., 1969, 1970) and Goddard Space Flight Center (Heppner, 1971), is displayed in **Figure 2** (from Haerendel, 1972). It exhibits the eastward and westward convection at high magnetic latitudes separating at a local time near 22:00. On the polar cap the sense appears to be reversed. Separating the flow vectors according to geomagnetic conditions showed

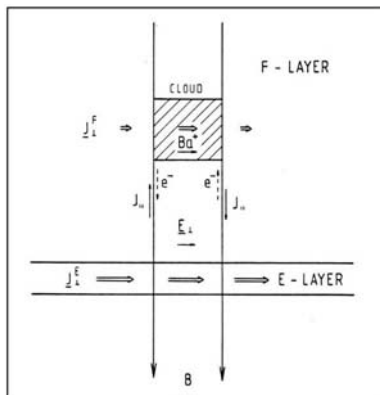


FIGURE 1 | Interaction of a finite cylindrical cloud with the lower ionosphere. Horizontal arrows, $\vec{J}_{1,2}^E$, are ion currents, the suffix 2 referring to field-lines passing through the cloud. Vertical arrows, $\vec{J}_s$, are field-aligned electron currents, and $\vec{E}_{1,2}$ the electric field. The cartoon shows that part of the E region Pedersen current is rerouted through the barium cloud by exchanging field-aligned currents. This leads to a reduction of the polarization field caused by the enhanced total conductivity and thus affects the observable motion of the cloud (Haerendel et al., 1967).

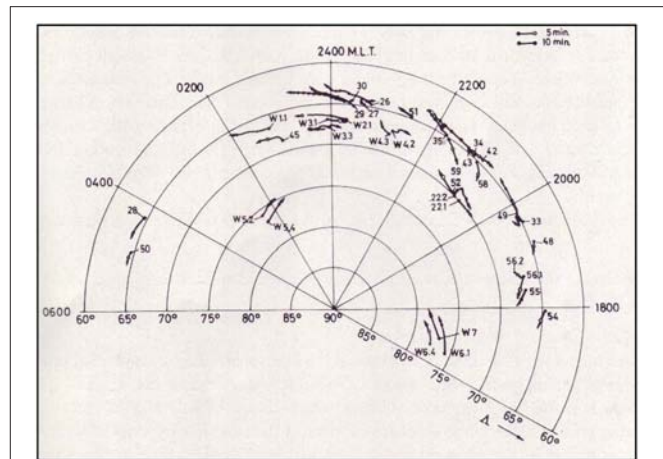


FIGURE 2 | Summary of the paths of barium clouds projected into an invariant latitude-magnetic local time plane. Open circles end a 5 min interval, full circles a 10 min interval (Haerendel, 1972). Paths labeled W refer to data from Wescott et al. (1969, 1970) and Heppner (1971). It exhibits the eastward and westward convection at high magnetic latitudes separating at a local time near 22:00 and a reversed sense of convection on the polar cap.

already the typical ordering which was to be fully elaborated in subsequent years, both by interpreting ground magnetic perturbations in terms of overhead Hall currents (e.g., Heppner, 1969) and directly from electric field measurements by double probes (e.g., Mozer and Bruston, 1967; Aggson, 1969). Not being restricted to twilight conditions and a few launch sites, the latter methods made these products of chemical release experiments on the long run obsolete. However, simultaneous flights of double probes and plasma cloud experiments proved very useful for intercalibration of the two methods (Fahleson et al., 1971; Kelley et al., 1975).

Barium cloud experiments in the ionosphere had quickly become very popular. Besides the above mentioned groups also the Air Force Cambridge Research Laboratories (AFGRL) (Rosenberg, 1971) took up this technique already in 1967. Experiments were not only performed at high magnetic latitudes but also at mid-latitudes and near the magnetic equator. Useful insights were obtained, as for instance into the driver of the Sq-current system. A collection of data from releases at mid-latitudes provided electric field measurements supporting the dynamo theory of Stewart (1882) with the neutral wind as driver (Haerendel and Lüster, 1968a). Releases from Thumba/India confirmed the peculiar behavior of the equatorial ionosphere at sunset and sunrise as inferred from incoherent scatter data. We will return to this topic in section Modification Experiments.

There are, however, situations in which plasma cloud experiments can yield unique insights. One of them is the motions and distortions of the seeded plasma in the neighborhood of auroral arcs. Furthermore, propagation or expansion parallel to the magnetic field can at times yield valuable information on the magnetic field direction and its variation, and most importantly, on parallel acceleration. A few results will be discussed.

Relative Motions of Plasma Clouds and Auroral Arcs

A fascinating variety of situations has been met when barium plasma clouds were injected in the neighborhood of auroral arcs. They ranged between parallel alignments of the cloud paths with the orientation of the arcs and greatly different relative motions, including crossings of the two. The first type of observations was made during magnetic quiet conditions and revealed that the arcs were embedded in the general plasma convection. The electric field was typically found to point normal to the arcs (Wescott et al., 1969). The second type of findings was obtained during strong magnetic activity. It transpired that the source of energy injected into the arcs was not frozen into the magnetospheric plasma but originated from intrusions of energy arriving from progressively further poleward.

Figure 3 documents a striking example of the latter. The event resulted from the injection of a barium jet along $\vec{B}$ from 540 km altitude by means of a shaped charge. It was performed in the dusk sector during an ongoing substorm further east (Wescott et al., 1975). The Ba^+ jet traveling upward to high altitudes had already split into several east-west separated streaks, when the aurora south of them was activated, forming a spiral structure and propagating poleward opposite to the equatorward drift of barium streaks. Crossing by the aurora had no effect on the plasma convection as manifested by the barium plasma. Another example occurred during a substorm about 3 h before magnetic midnight (Kelley et al., 1975). It was astounding how fast a rather irregular arc propagated poleward across the southeasterly drifting barium clouds. Both cases demonstrated the same facts, namely that the influx of new energy into the magnetosphere during substorms (and possibly also at other occasions) proceeded at progressively further poleward located flux tubes. However, it must be noted that, at the time of

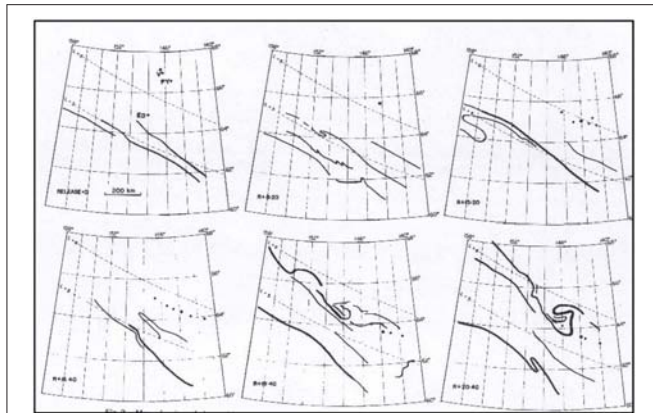


FIGURE 3 | Maps showing relative positions of the aurora and the 100 km projected positions of individual barium streaks. The initial release locations are at Venetie (V), Ester Dome (ED), and Fort Yukon (FY). After 16 min a substorm starts with new arcs expanding poleward and intercepting the equatorward moving barium streaks [Reproduced from (Wescott et al., 1975) with permission of John Wiley and Sons].

these experiments, the understanding of the connection to events at the outer boundary of the magnetosphere was at best rudimentarily developed.

The distribution of the energy injected into the magnetosphere from the tail is accompanied by different types of aurora. Even today not all of them are understood. The finding that some Ba^+ clouds were moving more-or-less parallel to an arc was not confirmed by other experiments. Often path or elongation axis of the Ba^+ clouds formed substantial angles with the auroral arcs (e.g., Haerendel et al., 1969). Later investigations substituting plasma clouds by incoherent scatter measurements of plasma drifts revealed the quite normal existence of proper motions between arcs and the environmental plasma (Haerendel et al., 1993). Such proper motions demonstrate that maintenance of the energy supply requires propagation of the arc into the energy reservoir constituted by sheared magnetic fields. This connection had been first proposed by the author in 1980 and subsequently elaborated in several papers (see Haerendel, 2007).

Field-Aligned Current in the Outer Magnetosphere

In 1969, the European Space Research Organization ESRO launched its first highly eccentric satellite, HEOS 1. It carried among others a barium release canister. It was ignited in March 1969 at a distance of 12. R_E (Haerendel and Mende, 2012). Closer look at the orientation of the long axis, obtained by triangulation from widely separated ground stations, showed a significant variation during the half-hour visibility. Further look into the onboard magnetometer data and inspection of the ground magnetic perturbations in the foot area in the neighborhood of Godhavn/Greenland revealed in-phase variations in the ionosphere. This was the first simultaneous measurement of field-aligned electric currents between the outer magnetosphere

and the ionosphere (Haerendel et al., 1971), at a time when one had just begun to find signatures of $j_{||}$ from magnetic field measurements at low-orbiting satellites (Zmuda et al., 1966, 1970). This example just exhibits one of the pioneering aspects of the plasma injection experiments in space with, of course, little sustainable value.

A later example of the above was the observation by Wescott et al. (1975) that the orientation of barium streaks produced by shaped-charge injections deviated appreciably from the magnetic field model. This was attributed to strong field aligned sheet currents flowing nearby and even allowed determination of their magnitude.

Field-Parallel Acceleration

The essence of auroral arcs is the conversion of free magnetic energy into kinetic and thermal energy of accelerated electrons and ions. Theories of the related existence of parallel potential drops were developed starting in the late 60's (Block, 1972; Swift et al., 1976) based on measurements of strongly field-aligned auroral electron distributions above auroral arcs (McIlwain, 1960; Evans, 1974). It was therefore most desirable to somehow succeed injecting Ba^+ ions into an auroral acceleration region. Not only did it require long waiting periods before conditions suited for a barium injection experiment close to an aurora were met, but hitting an acceleration region required multiple tries and was more or less a matter of luck. Since the acceleration regions extended from about 2,000 km up to and above 8,000 km, Brunner et al. (1970) from MPE and Wescott et al. (1972) from the University of Alaska independently developed the shaped charge injection of barium jets propagating with peak velocities of about 14 km/s and thus capable of probing heights of up to 30 000 km. An experiment with this technique from Søndre Strømfjord/Greenland in January 1975 turned out to be a lucky occasion (Haerendel et al., 1976). It was aided by the intrinsic proper motions of the aurora relative to the background and the injected barium jet. Initially located in an area of scattered auroral arcs, it expanded upward according to the adiabatic motion of ions. After 12 min an auroral arc appeared in the neighborhood of the jet. The low-velocity part of the jet, which meanwhile had sedimented to an altitude of 260 km, showed strongly enhanced transverse electric fields. For a short while, the jet was lost from observation because of the decreasing brightness of the strongly elongated streak. It could be recovered, but only by the most sensitive TV camera at the Thule observing site. Lacking proper triangulation the height distribution was determined by fitting with model field lines. However, distortion of the field by neighboring field-aligned currents introduced substantial uncertainties which are indicated in the data presented in **Figure 4**. In any case, a gain of energy by several keV is clearly indicated. Interestingly, the other initially separated streak did not exhibit any acceleration. This shows that auroral acceleration is restricted to narrow current sheets.

Heppner et al. (1981) found another way for barium clouds to reach high altitudes. They were generated from a satellite orbiting at 965 km altitude. By the diamagnetic force, $-\mu \nabla B$, the orbital momentum imparted to the ions was converted into parallel momentum in the upward diverging magnetic field. From the

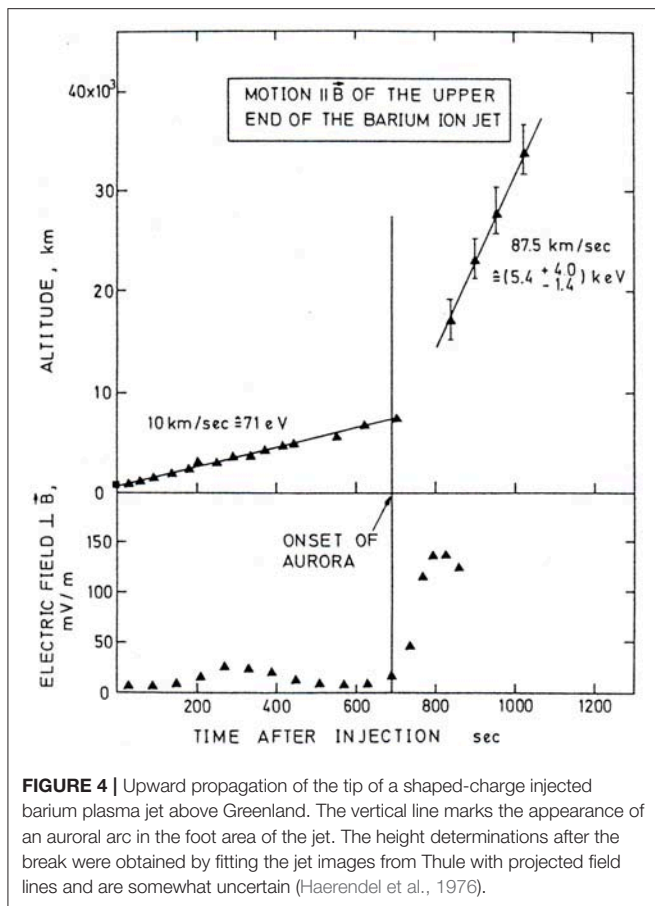


FIGURE 4 | Upward propagation of the tip of a shaped-charge injected barium plasma jet above Greenland. The vertical line marks the appearance of an auroral arc in the foot area of the jet. The height determinations after the break were obtained by fitting the jet images from Thule with projected field lines and are somewhat uncertain (Haerendel et al., 1976).

analysis of the field-aligned motions minor accelerations and decelerations were derived, but in one case substantial gain of parallel energy by 6 keV appeared to have occurred at heights above 15,800 km. The specific virtue of the releases from an orbiting spacecraft was that, contrary to shaped charge injections, they allowed measurements of very weak $E_{\parallel}$ fields and separation of spatial and temporal variations.

An intimately related result pertaining to the existence of field-parallel potential drops was the splitting of a streak at altitudes above 7,000 km by Wescott et al. (1976). It will be discussed in the next section in the context of coupling to the ionosphere. Taken together the few successful manifestations of field-parallel potential drops above auroral arcs can only be regarded as a proof of principle. No further insights into the acceleration process as such could be extracted. These were later on obtained from electric field and particle diagnostics from low-orbiting satellites, most importantly from the FAST mission (Carlson et al., 1998). Indeed, practically all aspects of the experimentation with plasma clouds discussed in this section were soon after successfully explored and even routinely measured by *in-situ* particle and field measurements. All the same, the barium plasma cloud experiments had the charm of being first, simple, visible, producing trustworthy results, and to fascinate professional or accidental observers. Already

in 1968 this was honored by an article in Scientific American (Haerendel and Lüst, 1968b).

COUPLING TO THE IONOSPHERE

Distortions and Striations

Barium plasma clouds are fascinating objects because of the visible manifestation of a variety of intriguing physical processes. A particularly striking aspect of experiments in the auroral ionosphere is the often observed strong distortions of the clouds. **Figure 5A** shows clouds over Ft. Churchill in 1967, one just being born out of the neutral barium gas, the other one generated 3 min earlier already extended transverse to the field by more than 200 km. The distortions of the gross shapes and the ubiquitous appearance of striations along the magnetic field direction, as shown in **Figure 5B**, have triggered much theoretical interest. Both experiments were performed in the auroral ionosphere where the electric fields are usually high. Distortions and striations result from interactions with the lower ionosphere, where the transverse electrical conductivity maximizes. As sketched in **Figure 1**, current exchange between cloud level and E-region by field-aligned electron currents can short-circuit or at least reduce electric polarization fields caused by conductivity enhancements by the Ba^{+} ions. This has the consequence of changing the conductivity in the E-region due to the divergences of the transverse Pedersen current at the interfaces with the field-aligned currents. The thereby generated secondary polarization fields act back on the motion of the barium ions. The artificial plasma clouds have mostly been injected at heights above 250 km. At these levels the ions are dominantly subject to $E \times B$ drifts. Any gross modifications of $E_{\perp}$ are therefore translated into shear flows, i.e., into distortions.

Equally striking are smaller-scale perturbations growing into visible striations. In the rear of a moving cloud, density (i.e., conductivity) enhancements are slowed down in comparison with the undisturbed motion because of the reduction of $E_{\perp}$. They lag behind and become more pronounced. Density depletions, on the other hand, experience a higher electric fields and can advance faster into the cloud. Thus, the rear side of a drifting cloud is unstable with respect to small-scale perturbations. The opposite holds for the front side.

What has just been described has soon been recognized by Linson and Workman (1970) as leading to a cross-field or gradient drift instability. The non-linearity of the cloud-ionosphere interactions asked for treatment by numerical simulations. The stage was set by the two companion papers of Perkins et al. (1973) and Zabusky et al. (1973) and independently by Lloyd and Haerendel (1973). Indeed observed bifurcations (**Figure 6**) or splitting into multiple structures on the rear side of the distorted cloud were successfully reproduced by McDonald et al. (1981). In that paper the further question is pursued why structures of about 1 km width appear to be “freezing up,” i.e., persist for long times, while the gradient drift instability should lead to splitting into smaller-scale structures. Rocket flights through barium clouds measuring the electron density indeed verified the existence of large-amplitude perturbations with scales well below 1 km matching a power law distribution with a

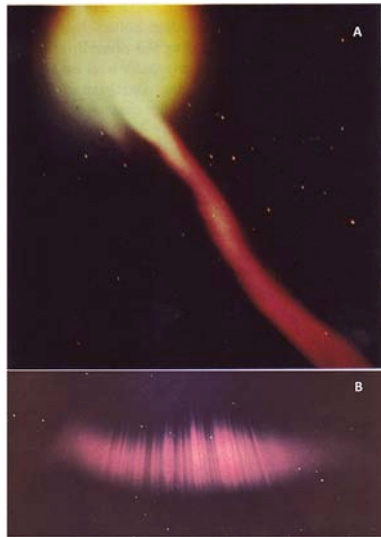


FIGURE 5 | (A) Two barium clouds injected above Ft. Churchill/Canada with 3 min separation in August 1967. The yellow/blue colors indicate neutral Ba, the purple color the first injected cloud, strongly extended over 200 km horizontally. **(B)** Highly striated and extended barium cloud, also over Ft. Churchill (Haerendel and Lüst, 1968b).



FIGURE 6 | Barium cloud over Kiruna (1968) viewed along the magnetic field (Haerendel, 1996).

spectral index between -2 and -3 (Baker and Ulwick, 1978). McDonald et al. (1981) attributed the stabilization of km-scales to the presence of a hypothetical turbulent diffusivity which, depending on strength, could stabilize km-scale structures.

An analytical approach to the striation problem was chosen by Völk and Haerendel (1971). They took into account the feedback on density perturbations developing in the cloud from the density changes arising in the background plasma from the short-circuiting Pedersen currents. Because of a phase shift of the polarization field by 90 degrees between the density variations in the cloud and those in the background, the image striations in the background are tilted with respect to B . Furthermore, the finite parallel conductivity was included in the stability analysis. This has the consequence that the effective integrated Pedersen conductivity is reduced below a certain transverse scale and with it the depolarization effect. As a result the growth rate of a density perturbation or striation decreases strongly below a certain thickness.

In summary we have two competing explanations of the “freezing-up” of scales above about one km, turbulent diffusivity or reduced depolarization. The next sub-section introduces another but related aspect.

Magnetosphere-Ionosphere Coupling

With few exceptions barium cloud experiments were performed in or, by the shaped-charge technique, out of the ionosphere. A barium cloud injected from a Scout rocket at a geocentric distance of $6 R_E$ in 1971 presented a new aspect, a prolonged phase of acceleration toward the flow speed of the ambient magnetospheric plasma. **Figure 7** shows the development of the cloud during 1 h of observation. A few minutes after injection the cloud was seen to split into several magnetic field-aligned structures. While the brightest streak, the maximum-density core, followed closely the rocket trajectory, the other streaks with gradually lower brightness (density) separated with increasingly higher speed. In the first report on that experiment at a COSPAR conference (Haerendel, 1973) only preliminary results were presented with scarce speculative interpretations. Coupling to the ambient plasma and steady momentum deposition in the ionosphere was recognized as the most intriguing aspect raised by the experiment, but no detailed analysis was attempted. The greatest surprise was the apparent absence of oscillatory motions with a period of the bounce time of an Alfvén wave between cloud and ionosphere as predicted by the seminal theory of Scholer (1970). The rather different speeds assumed by the structures were (falsely) related to local gradients and variable ambient flow speeds.

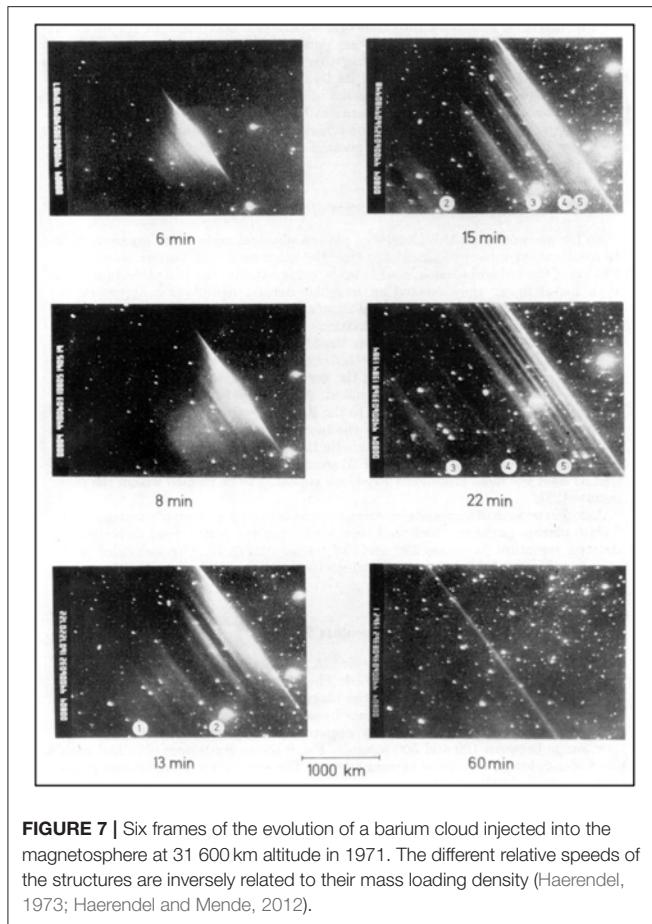
The other striking observation, made by the onboard search coil magnetometer, was the formation of a diamagnetic core. Entry into the returning field after one half minute indicated a width of the boundary much thinner than the gyroradius of a Ba ion. The true nature of the magnetic field return transpired much later at the occasion of the artificial comet experiments (s. section Interaction With the Solar Wind).

In retrospect it is hardly understandable that the evaluation of this most challenging experiment was left in this rudimentary stage. This can only be attributed to the lack of manpower and the pressure exerted by the close sequence of campaigns in various remote places, but not be excused. Publication in a refereed journal had to wait for 40 years. Finally, the treasures inherent in that experiment were finally unearthed.

Haerendel and Mende (2012) first reviewed the experimental parameters of cloud and environment and, following (Scholer, 1970) determined the magnitude of the theoretical coupling time constant to the environments:

$$\tau_0 = \frac{\mu_0 \rho_c \ell_{||} v_A}{2 B^2} \quad (3)$$

With ρ_c being the mass density in the respective streak of the cloud, $\ell_{||}$ the effective length of the streak, v_A the Alfvén speed, and B the ambient magnetic field strength. It turned out to be almost 4 orders of magnitude larger than the bounce period of an Alfvén wave. With the high reflectivity of the ionosphere, little momentum would be transferred to



the ionosphere but instead stored in increases of the magnetic shear stresses. Application of Scholer's theory predicted that the relative velocity of the cloud should have decreased by an order of magnitude after about 15 bounce periods (5 min) or even been reversed.

None of this was observed, but rather constant accelerations of the streaks, the greater, the lower the respective density/brightness. This meant that, after a short initial short interaction with the ionosphere by exchanging Alfvén waves, each streak had found its coupling constant in accordance with its respective mass content. How was that achieved? After a few reflections, observationally not resolvable, the streaks had assumed widths with scales below 1 km, when mapped into the ionosphere. Scholer (1970) had already suggested that owing to the finite parallel resistivity, $\sigma_{||}$, the integrated Pedersen conductivity would have to be replaced by a substantially smaller effective Pedersen conductivity, $\Sigma_{P,eff}$. The potential drop inside the ionosphere would not be controlled solely by the Pedersen conductivity, σ_p , but by a geometric mean conductivity, $\sigma_m = \sqrt{\sigma_{||} \cdot \sigma_p}$. This would reduce the reflectivity of Alfvén waves and thus lower the braking time.

Haerendel and Mende (2012) went one step further and postulated that scale breaking of the initial cloud was the reason for the formation of the striations. Perfect matching between wave impedance and effective conductivity, i.e.,

$R_w \cdot \Sigma_{P,eff} = 1$, yields zero reflectivity. In this case the coupling constant is:

$$\tau_o = \frac{\rho_c \ell_{||}}{2 B^2 \Sigma_{P,eff}} \quad (4)$$

The principle behind scale-breaking is the optimization of the energy dumping into the ionosphere (see Haerendel, 2014). This energy resides in the magnetic tensions acquired by the initial reflections at the ionosphere. The consequence of perfect matching is that the transverse electric field of the striation is applied to the ionosphere and decays inside of it. The existence of a substantial potential drop above the ionosphere was discarded by the authors. Thus, the barium structures moved decoupled from the frame of the ambient plasma but were continuously losing momentum and kinetic energy with respect to the plasma frame until reaching the ambient flow speed. The leading diffuse streak in Figure 7 was indeed seen to acquire a velocity consistent with the external flow speed.

It is interesting to compare the here discussed coupling process with that of the ionospheric releases. The difference lies in the ratio of the time scales for momentum coupling with growth times of any deformation. In case of magnetospheric plasma injections, the ratio is large. The opposite holds for ionospheric clouds. Effective depolarization, as discussed in the preceding sub-section, means effective coupling. Formation of smaller-scale striations is enabled by decoupling due to reduction of the effective Pedersen conductivity thus allowing instabilities and relative motions. An interesting intermediate case was the Cameo releases from an orbiting spacecraft at 965 km, already discussed in sub-section Field-Parallel Acceleration (Heppner et al., 1981). All releases formed sheets of series of field-aligned streaks, obviously owed to the above discussed scale-breaking process. It would be rewarding if one could re-evaluate the observed irregularities of the initial transverse motions in terms of a finite momentum coupling time.

It is interesting to compare the above described experiment with later releases at various heights in the magnetosphere by the Combined Release and Radiation Effects Satellite (CRRES) (Fuselier et al., 1994). Depending on geocentric distance (2.0 R_E , 3.4 R_E , 4.8 R_E), i. e. on magnetic field strength and plasma density, the coupling times to the frame of the ambient medium varied from 25 s to 5 min. This was entirely attributed to the momentum coupling to the ambient plasma. Only in one of the releases (G-3) there were indications of reflection at the ionosphere. A comprehensive overview of the CRRES experiments has been provided by Bernhardt (1992).

Decoupling by Potential Drops in the Magnetosphere

Decoupling from the ionosphere by processes within the magnetosphere can occur basically in two ways, by field-parallel potential drops in field-aligned current sheets or within kinetic Alfvén waves, owing to a finite $E_{||}$ component. In the first case, the macro-physical reason is the mirror effect on the low-density energetic electrons carrying the current. The micro-physical

reasons for sustaining the parallel electric fields are smaller-scale turbulence and, in particular, the action of solitary waves. While quasi-stationary auroral current sheets have typically widths of the order of 10 km, propagating kinetic Alfvén waves have considerably smaller transverse scale comparable to the electron gyroradius. What would we expect from barium plasmas injected into such acceleration region? Two effects, upward acceleration of the positive ions accompanied by decreasing brightness and decoupling from magnetospheric convection as indicated by accompanying lower-altitude clouds. The ion jet experiment from Greenland, discussed in sub-section Field-Parallel Acceleration, showed upward acceleration, but the decoupling, although very likely, could not be confirmed because of the lack of proper triangulation of the upper streak.

An experiment by Wescott et al. (1976) exhibited a reversed behavior. Clear decoupling was seen but very little upward acceleration. Barium ion jets were injected by rocket flights from Alaska into a diffuse aurora in the dawn sector. **Figure 8** shows the development of the streak. At 15 min after injection a brightening of almost a factor of 3 was observed above an altitude of 5,500 km accompanied by a splitting into two secondary rays above that altitude. This has been interpreted as a decoupling by a parallel potential drop. The origin of the brightening was attributed to a Doppler shift out of the minimum of the Ba^+ absorption line at 4,554 Å, caused by a sudden increase of the combined upward and transverse ion velocities away from the sun. Quantitative checks supported this explanation. Only relatively low velocity increases were needed. And curiously, evaluation of the upward motions suggested an energy gain of only 34 eV, well below what is typical for auroral arcs.

The reason for the absence of any strong upward acceleration observed probably lied in the auroral situation. The barium streak was determined to nearly coincide with the poleward edge of an omega band. At this boundary exists a strong flow shear from mildly eastward inside the visible bands to a rapid eastward speed of up to 4 km/s (Buchert et al., 1988). Furthermore, from fitting the high-altitude main streak with magnetic field models, Wescott et al. (1976) derived the existence of an upward sheet current of 0.2 A/m. This is in good agreement with the nature of omega bands which are caused by precipitation of typically 10 keV magnetospheric electrons carrying rather weak parallel currents of order $1 \mu\text{A}/\text{m}^2$ distributed over 100 km north-south or more (Buchert et al., 1990). This means that for sustaining the current no post-acceleration of the electrons is needed. There is a dilemma. The observed splitting of flux tubes filled with barium plasma above 5,000 to 6,000 km is not disputable and no other obvious reason for that than electrostatic decoupling comes to mind. On the other hand, the geophysical nature of the local aurora does not require parallel potential drops, nor do the Ba^+ ions indicate their existence except for a few tens of volt. Can it be that there are indeed weak potential drops existent at the interface between the dilute hot magnetosphere and the cooler and denser topside ionosphere? The one-dimensional model of electric fields in upward field-aligned currents by Ergun (2002) showed exactly that possibility. At the time of writing of this review, this appears to be the best explanation for the discussed observations.

The here presented evidence for magnetospheric decoupling regarded together with the prior discussion of coupling and decoupling, including the evidence for field-parallel acceleration in Sub-section Field-Parallel Acceleration, represent examples of the unique contributions of the plasma cloud experiments to the exploration of the fascinating subject of magnetosphere-ionosphere coupling.

Interaction With the Solar Wind

The proposal of Biermann et al. (1961) to generate plasma conditions in the solar wind simulating its interaction with comets, was the first goal of German space research. It took a quarter century, space around Earth was already widely explored, before this goal was fulfilled. Only two such “artificial comet” experiments were performed within the AMPTE mission (Krimigis et al., 1982). However, the presence of extended *in-situ* diagnostics in support of the optical tracing from ground and aircraft led to a rich harvest of new physical insights. AMPTE was standing for Active Magnetospheric Particle Tracer Explorers and consisted of three spacecraft from the USA, Germany, and the United Kingdom. The name was derived from the other goal of the mission, namely to investigate the mass transfer efficiency from solar wind and magnetotail into the inner magnetosphere by injecting primarily Lithium ions from the outer spacecraft, the Ion Release module (IRM), and searching for them with the Charge Composition Explorer (CCE) inside the magnetosphere. Unfortunately, this part of the mission was not successful (Krimigis et al., 1986). Here we will only summarize the physical processes taking place at the location of the injected plasma.

Pick-Up Ions

The contribution of interstellar helium to the ionic constituents of the solar wind by photo-ionization of the neutral atoms near the sun had been a subject of investigation long before the actual pick-up of freshly ionized lithium ions was recorded by the AMPTE/IRM spacecraft. A first Li-cloud was injected into the solar wind ahead of the bow shock for the tracing purposes mentioned above in September 1984. Being ionized with a time scale of 1 h, Li^+ ions were continuously generated in the freely expanding neutral cloud and became immediately subject to the force of local electric field. Initially they just followed the electric vector, but when born at larger distance from the spacecraft and thus sufficiently accelerated, the Lorentz force took over. As seen by the plasma analyzer (Paschmann et al., 1986) and the Supra-thermal Energy Ionic Charge Analyzer (SULEICA) (Möbius et al., 1986), the energy of the ions grew with time and the arrival direction changed. Only weeks after the latter instrument registered for the first time interstellar He-ions *in situ* (Möbius et al., 1985). This was the beginning of a long success story of interstellar pick-up ions and their role in the solar wind (Möbius et al., 1988). The AMPTE mission was first in actually observing the very pick-up process.

Magnetic Cavities

A genuine plasma physical effect of the Li^+ injections was the formation of magnetic cavities. Even with the long ionization

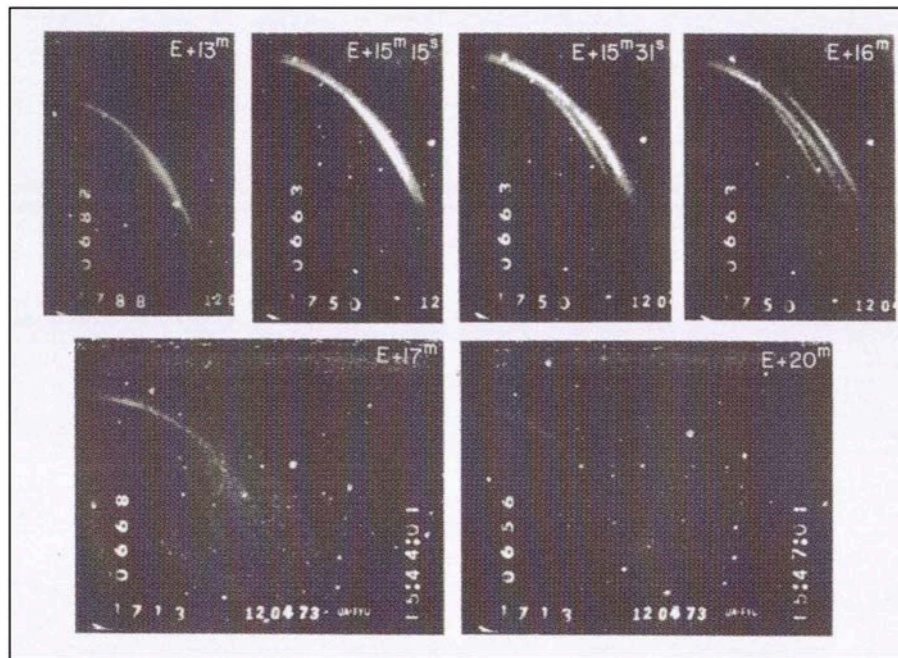


FIGURE 8 | Plasma jet injected from Alaska showing anomalous brightening after 15 min and 15 s above an altitude of 5,500 km accompanied by splitting into two secondary streaks, which are decoupled from the lower sections of the respective flux tubes. Non-parallelism of the streaks and deviations from model field indicate nearby upward sheet currents (Reproduced from Wescott et al., 1976 with permission of John Wiley and Sons).

time of lithium and the more so with the fast ionization of barium, the freshly generated plasma caused a strong inflation and depression of the magnetic field, lasting until the momentum of the plasma cloud equaled the external magnetic pressure. This was already seen in the Scout experiment in the magnetosphere (section Magnetosphere-Ionosphere Coupling). The altogether eight AMPTE releases offered a valuable sample of that effect allowing deeper investigation of the ongoing physical processes. The most important and hitherto unknown ones were the unexpectedly fast recovery of the cavity, the mass pile-up behind the front of the returning field, and the thinness and longitudinal structure of the magnetic boundary.

Figure 9A documents the first two effects as recorded during the first artificial comet experiment upstream of the bow shock on 27 December 1984 (Gurnett et al., 1986). Immediately after ignition of the Ba charges the magnetic field strength dropped below the level of sensitivity. After 80 s, a sharp return of the field with greatly amplified amplitude was accompanied by an increase of density by a factor of about 5. From estimates of time and radius of maximum expansion one could determine that the speed of the returning front of the field was 4.5 km/s. Haerendel et al. (1986) interpreted what had happened as consequence of a snow plow effect. The returning field sweeps up the yet unmagnetized Ba plasma which piles up behind the front. The high mass load slows down the speed of the returning field to the thus modified Alfvén speed. Combining magnetic gradient and curvature forces yielded a value of a Alfvén speed consistent with the derived speed of the snow plow. In the second comet experiment of July 1985 and even better in a later experiment

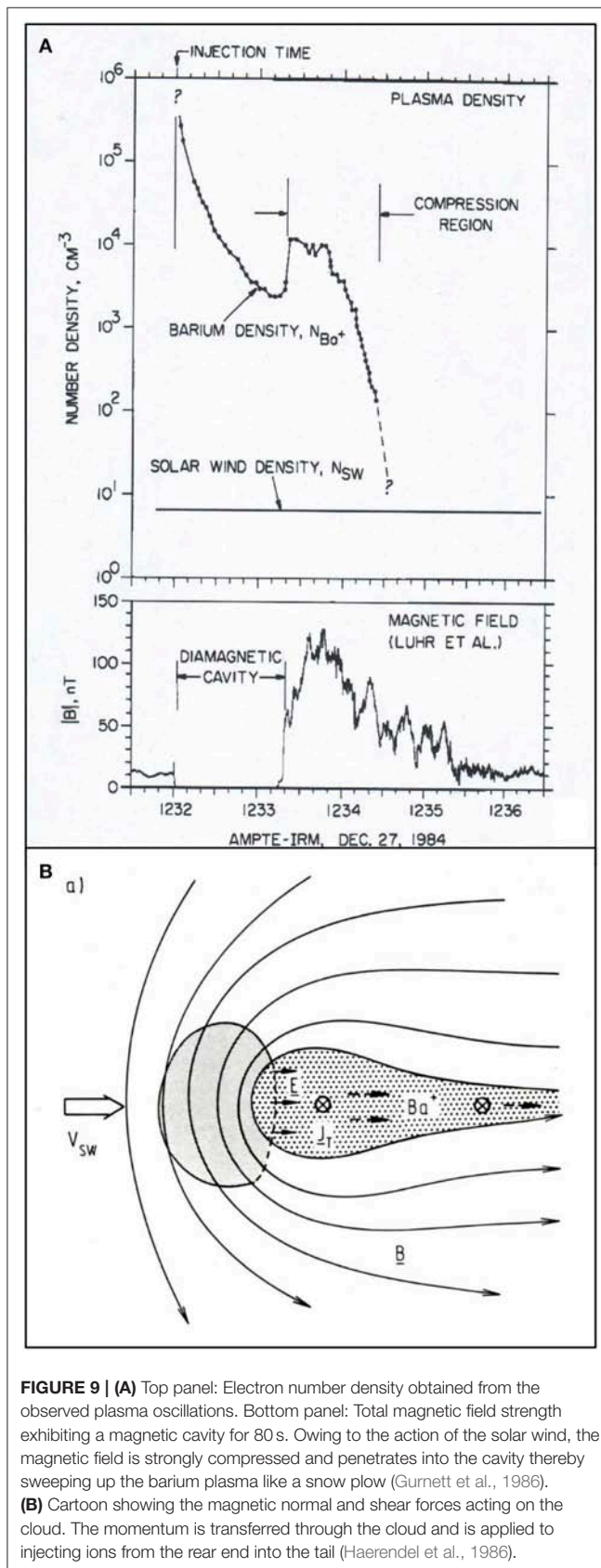
in the magnetotail the snow plow propagation could be clearly observed.

The magnetic field compression from 10 to 130 nT was attributed to the action of the solar wind. As sketched in **Figure 9B**, magnetic normal and shear stresses contribute about equally to the force acting on the plasma cloud. This leads to the compression factor (Haerendel et al., 1986; Haerendel, 1987):

$$\kappa = \frac{B_c}{B_0} = 2 M_{A\perp} = 2 \frac{v_{sw\perp}}{v_A} \quad (5)$$

With $M_{A\perp} \approx 6$ the relation can explain the observed amplification.

The returning magnetic field offered two surprises, the short duration of the entry of only 0.5 s and the low amplitude of the electric noise appearing below the electron gyrofrequency (Gurnett et al., 1986). Expected was a high noise level and broader width indicating the presence of some kind of anomalous magnetic diffusivity. The short duration of the electric noise and its low amplitude suggest the existence of an electron-scale sub-layer, through which the electrons are scattered into the snow plow front, which is substantially wider. This became clearer with a barium injection in the tail in the following year. In that experiment, the magnetic cavity and the propagating snow plow were observed more or less along the magnetic field and well resolved over many minutes. Most remarkable were the field-aligned ripples covering the inner surface of the already magnetized plasma (**Figure 10a**) (Bernhardt et al., 1987). This modulation was subsequently interpreted by the author in a



paper dedicated to Ludwig Biermann (in German) as owed to the different ways ions and electrons enter the magnetic field (Haerendel, 1986). As sketched in **Figure 10b**, the boundary is so thin that the ions are hardly deflected by the magnetic field but stopped by an opposing polarization field. The electrons are scattered into the boundary layer and, being magnetized, carry a Hall current shielding the magnetic field. Balancing it with the current carried by the entering ions, a width of the order of c/ω_{pi} was derived (Szegő et al., 2000, p. 617). The observed ripples had widths of about 300 km, whereas the ion inertial radius of the ions was about twice as large. By contrast, Bernhardt et al. (1987) attributed the ripples to an interchange instability, but obtained scales at least one order of magnitude too low. There have been other interpretations (see below).

Many years later the author noticed the similarity between barium ions entering the returning magnetic field and a situation encountered in the central plasma sheet during the onset of substorms (Haerendel, 2015). He postulated the formation of a sharp boundary of width, c/ω_{pi} , where ions in the high-beta central current sheet are being slowed down without being reflected. The question, not asked in the barium experiment, is: Where do the energy and momentum go that are deposited by the ions? The answer came from the reverse side by wondering about the source of the energy of the highly structured thin arcs observed at substorm onset. The short-lived erratic bead structure was attributed to the impact of kinetic Alfvén waves launched from corresponding ripples of the Hall current in what was called the “stop layer.” In these ripples the kinetic energy of the ions is converted into electromagnetic energy and carried away by the waves. Since their amplitude is limited by the strength of the stopping magnetic field and Hall current, energy and momentum can only be deposited within a limited lateral scale of the order of c/ω_{pi} . While the stop layer in the equatorial magnetosphere at typically $8 R_E$ is still a conjecture and the ripples have not yet been observed, the barium experiment demonstrates their existence in a not dissimilar situation.

Magnetic cavities were also studied with releases of the CRRES mission (Huba et al., 1992). The emphasis was mainly directed toward comparison with numerical simulations by MHD and non-ideal Hall MHD codes. Structures of scales size 1–2 km formed within 3 s for the release at $4.8 R_E$ and with scale sizes of 10–15 km within 22 s for the release at $6.2 R_E$. They are attributed to the collisionless Rayleigh-Taylor instability. This differs decisively from the stop layer mechanism discussed above. Equally different is the reason for the density increase after the maximum expansion time in these simulations. It is a pile-up resulting from sequential deceleration, first of the fastest ions and followed by the slower ones which are catching up. Owing to the use of MHD a snow plow effect does not exist. Unfortunately, there is no accompanying documentation of the actual optical observations.

Ion Extraction and Momentum Balance

Among the many surprises of the AMPTE barium releases in the solar wind the perhaps most perplexing one at first sight was

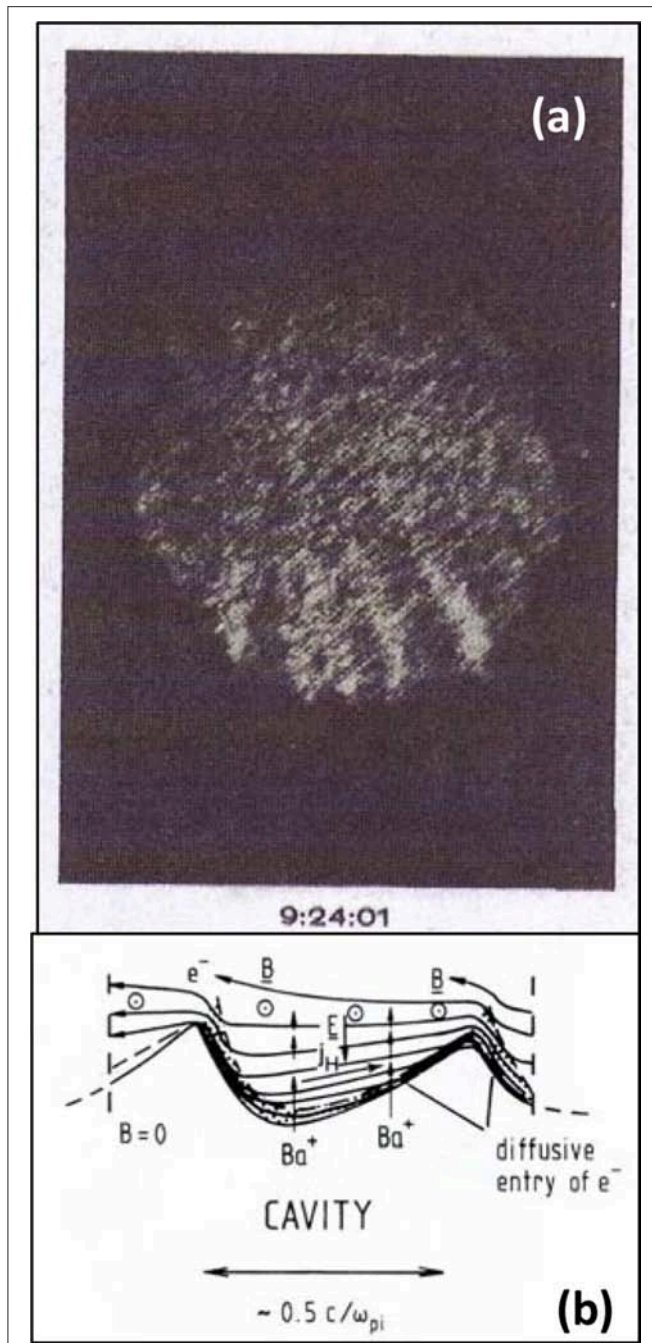


FIGURE 10 | (a) Ripples on a barium plasma cloud released in the Earth's magnetotail (Bernhardt et al., 1987). **(b)** Proposed formation of the ripples as caused by the ions penetrating directly into the field with the electrons following by diffusive entry. The electric field, E , stopping the ions drives a Hall current which is shielding the magnetic field. Balancing the energy entered by the ions with the energy carried away along the magnetic field (normal to the plane) enforces structuring of order c/ω_{pi} (Szegő et al., 2000).

the dynamic behavior of the comet head. During the first 4 min, the center of mass was not at all displaced in the direction of the oncoming solar wind but at right angles (Valenzuela et al., 1986). The explanation was given by Haerendel et al. (1986) and

in more detail in Haerendel (1987). In the frame of the barium cloud (comet head), there is an electric field, $E_{\perp} = -v_{sw} \times B$, which is felt by the ions in the boundary layer of the cloud. Ions are being extracted and enter a cycloidal path much wider than the cloud dimension. Momentum balance, i.e., the recoil of the extracting force, pushes the bulk plasma sideways, opposite to $E_{\perp}$. The long station-keeping of the cloud in s.w. direction is a consequence of the generation of a tail. The normal momentum imparted by the solar wind is transported through the comet head and delivered to its rear. From there ions are extracted, again by an electric force, and injected downstream. A cloud image from the first artificial comet experiment and an explanatory sketch are contained in **Figures 11a,b**.

Simple analytical calculations to be found in Haerendel (1987) and Szegő et al. (2000) lead to the following simple relations. The accelerating force experienced by the cloud via the draped and compressed magnetic field is:

$$g_{\parallel} \approx \frac{4\rho_0 u_{sw}^2 A_c}{M_c}. \quad (6)$$

A_c is the cross-section of the cloud, and M_c the total mass. In the factor 4 are contained the momentum imparted by the Alfvén wings of the draped magnetic field and the compression of the field (s. Equation 5). The transverse acceleration was determined to be:

$$g_{\perp} = \frac{1}{4} g_{\parallel} \quad (7)$$

The latter actually matches the observations of the lateral displacement (Valenzuela et al., 1986). Being applied to the tail formation, $g_{\parallel}$ could not be directly observed (s. below).

Figure 11b shows the flow lines of the electrons (left) and of the sw protons (right). The electrons are magnetized and follow the $\underline{E} \times \underline{B}$ drift, while the ions experience only mild deflections by the strong internal magnetic field. Both flows constitute Hall currents corresponding to the field concentration in the comet head. The magnetic field is asymmetrically distributed in the cloud owing to the different widths of the electron and ion Hall currents. It is the negative magnetic pressure gradient existing through much of the cloud what transfers the recoil of the ion extraction pushing the comet head sideways.

Because of the lateral deflection, the AMPTE/IRM exits the cloud in the direction of the ion extraction. When it had reached the low-density flank, suddenly (at 1234:23 UT) a very strong electrostatic noise appeared (Gurnett et al., 1986). The latter author interpreted the rather unstructured electrostatic noise as generated by an ion beam-plasma instability between the nearly stationary barium ions and the rapidly moving solar wind protons. In the frame of the $\underline{E} \times \underline{B}$ drifting electrons, the Ba-ions moving along the $\underline{E}$ -vector are practically at rest. Papadopoulos et al. (1987), expanded the theory along the same lines and argued that the noise was saturating by proton trapping in the electric wave field. From the spectral energy density in the lower hybrid region he determined an anomalous ion-ion collision frequency.

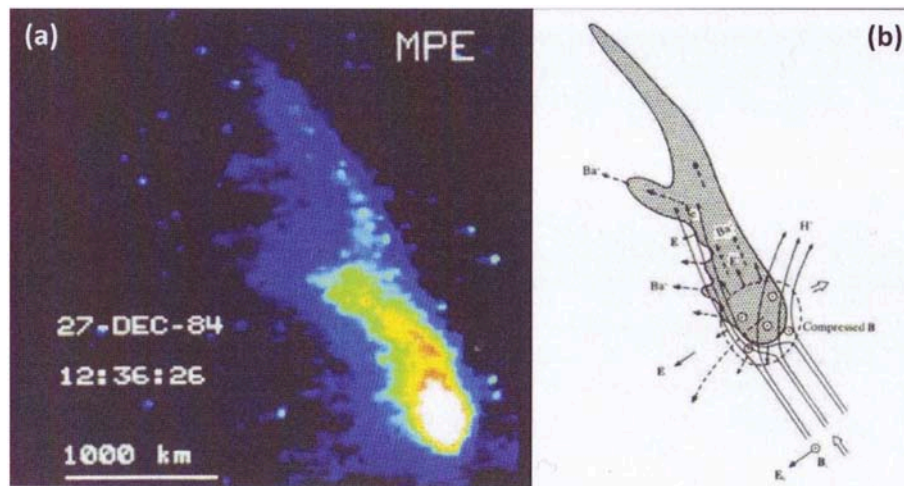


FIGURE 11 | (a) False color image of the first artificial (plasma) comet 4.4 min after release in December 1984 exhibiting ion extraction (to the left) and tail formation (Valenzuela et al., 1986). **(b)** Interpretation of the cause of the lateral displacement of the comet head (open arrow) as reaction to the extraction of ions by the solar wind electric field, E . Solar wind protons, H^+ , are only slightly deflected by the trapped magnetic field, whereas the magnetized electrons perform an $E \times B$ -drift along the opposite flank. The associated Hall current shields the field on one side (Haerendel et al., 1986).

This implies a dynamic momentum coupling from the protons to the barium ions in addition to the laminar momentum coupling by the extracting quasi-static electric field.

Tail Formation

As argued above, the station keeping of the cloud in the sw direction was owed to the momentum imparted by the draped magnetic field and transported through the cloud to its rear end. Like in the popular toy of several suspended balls in contact, it is the last mass that carries away the momentum. But how is that achieved in the Ba^+ cloud? A simple argument is presented in Szegő et al. (2000). In most of the cloud the upstream directed normal stress (magnetic and pressure) is balanced by the downstream pointing tangential stress of the field. The bulk plasma is not accelerated. Toward the rear, the relation changes. Both forces point downstream, only balanced by the inertial force of the accelerated ions. The acceleration completed, the density of the ion outflow must equal approximately that of the solar wind, since the ions are neutralized by the solar wind electrons entering the forming tail from the side and not by the photo-electrons (Haerendel, 1987). With that assumption the speed attained by the ions is (Szegő et al., 2000):

$$u_{||extr} \approx 2 \sqrt{\frac{2}{\mu_i}} v_{sw} \quad (8)$$

μ_i is the atomic weight of the ion. For Ba^+ this amounts to 24 % of the sw speed. Photometry of the plasma cloud was not really possible. However, the pictures obtained by the Doppler imaging system of Rees et al. (1986) show clearly the strong drop in density to be expected from the above considerations.

Computer Modeling

Naturally, the above presented simple analytical derivations are very simplistic, just trying to catch the essential physics. Reality is more complex. Fortunately the theoretical plasma physics community was fascinated by such plasma experiments not impeded by walls and analyzed on the microphysical scale by *in-situ* diagnostics. Many efforts have been made to study the physics by numerical simulations.

Cheng (1987) studied the lateral deflection of the cloud and emphasized the role of the deflection of the solar wind protons. It generates an electric polarization field driving a Hall current on the side opposite to the solar wind electric field. This maintains the compressed magnetic field and its normal stresses contribute to the sideways motion.

Brecht and Thomas (1987) performed a three-dimensional simulation with a hybrid code in which the electrons were considered as a massless fluid. They studied in particular the formation of magnetic cavities and the magnetic field compression within the snow plow region. They could also reproduce the cross-field deflection of the cloud.

Huba et al. (1987) applied theory and simulation of the Rayleigh-Taylor instability to the appearance of striations in the tail release, discussed and interpreted above. They see the driver in the slowdown of the initially expanding cloud by the outer magnetic field. In the large Larmor radius limit, the fastest growth rates were found at the shortest wavelengths. This is in agreement with the analysis of Bernhardt et al. (1987) (see above), but not consistent with the observation. However, the simulations also showed the longer wavelengths to dominate the forming structures.

Bingham et al. (1988) simulated the initial magnetic field compression around the barium cloud and addressed in addition the observed heating of the electrons. Calculating the beam-plasma instability and the growth rate of lower-hybrid waves

they found substantial electron heating by Landau damping along the magnetic field. However, the observed electron spectra, as measured by the United Kingdom Sub-satellite (UKS) only a bit outside the cavity (Bryant et al., 1985), were more alike to an acceleration by a quasi-static electric potential drop.

An interesting approach was the bi-ion fluid simulations of Sauer et al. (1994) and Bogdanov et al. (1996). Light and heavy ions are considered as fluids and the massless electrons follow from charge neutrality and Ampere's law. Christian Fischer in Szegő et al. (2000) extended the bi-ion fluid simulations to three dimensions and computed the case of a plasma comet. In **Figure 12** one can clearly recognize the different paths taken by the heavy ions and the solar wind protons. Furthermore, it reproduces the observation, more striking with the second experiment in 1985, that the tail flows are not laminar but form clumps or knots. The author had suggested a clumping instability for their formation (Szegő et al., 2000). The simulations confirm the underlying mechanism that the heavy ions move in and out of the knots and are thus on average faster than the knots are proceeding. This has implications on so-called disconnections in comet tails.

In summary of this incomplete account of the various processes observed in the two artificial comet experiments, one can say that they led to the recognition that fundamental transport processes are often enabled just by the different inertia of ions and electrons without much accompanying noise.

MODIFICATION EXPERIMENTS

Inospheric Modifications

There are many ways of injecting neutral gases into the upper atmosphere, both intentionally and as by-products of rocket exhaust. Once released, some gases, like barium, strontium, europium are photo-ionized or, in special cases, can be ionized by collisions. Others undergo chemical reactions by which the natural plasma component is reduced by recombination processes, thus creating plasma holes rather than clouds (Mendillo, 1988). The latter techniques and applications have been reviewed by Bernhardt et al. (2012). Here we summarize the attempts to modify the ionosphere for stimulating, triggering, or at least tracing natural ionospheric instabilities. The formation of striations in barium clouds discussed in section Coupling to the Ionosphere is an example.

A most ambitious goal was the attempt to trigger equatorial spread-F. This phenomenon has attracted much attention. It is an ionospheric realization of the Rayleigh-Taylor instability after sunset when the gradient of the denser F-region has risen up to and above 350 km. A wide range of secondary processes follows (Hudson and Kennel, 1975; Zalesak et al., 1982). It has been explored primarily by radar measurements, optical observations, and by *in-situ* probing with sounding rockets (Balsley et al., 1972; Woodman and La Hoz, 1976; Kelley et al., 1982; Kelley, 1989). MPE conducted four rocket campaigns under spread-F conditions, initially for tracing the evolution of the irregularities and finally for triggering an instability. The first three attempts took place in Thumba/India in 1972, in Natal/Brazil in 1973, and

in Punta Lobos/Peru in 1979. Either the rockets failed or spread-F did not develop, as long as the clouds were observable. The only result was the appearance of the ubiquitous striations and distortions. However, in the combined campaigns of BIME by Narcisi (1983) and Colored Bubbles by Haerendel et al. (1983) in Natal in 1982, there was partial success. Later theoretical considerations made clear that charges with 26 kg of the Ba-CuO mixture as employed in these experiments would not suffice to generate sufficiently heavy plasma clouds capable of massively disturbing the ionosphere. Much higher masses would be needed. However, they sufficed to excite localized spread-F.

The basic idea of the Colored Bubbles experiments was to place two big barium clouds side by side, on upleg and downleg of a rocket flight, respectively. It was to be expected, that while the two flux tubes heavily loaded with barium plasma would move downward, the flux tube in between would rise. This is owed to the incompressibility of the magnetic field. The hope was that eventually the rising flux tube would evolve into a bubble, the name introduced for a low-density plasma rising into the dense upper F-region. Such bubbles, specifically their westward borders, have been recognized as the main sites of secondary instabilities due to the eastward blowing neutral wind (Zalesak et al., 1982). However, numerical simulations of the experimental situation by Çakir et al. (1992) later showed that the creation of a real bubble would require 40 times higher masses to be injected.

Besides the two big barium releases near apogee, there were five small releases on the upleg, in order to trace the vertical shear of the plasma flows, and a light europium release for tracing the expected plasma uplift between the big clouds. Above about 250 km the flows were directed eastward. This meant that the barium clouds disappeared in the shadow after about 20 min. This precluded optical tracing of the longer-term development. For this reason the two campaigns had secured the cooperation of several groups with competence in radio wave tracing by ionosondes, radar, or recording scintillations of radio beacons from geostationary satellites. The latter was performed by Johnson and Hocutt (1984) on a NASA aircraft. Following the eastward motion of the clouds, they were able to track scintillations with steadily growing amplitudes consistent with the continued eastward drift of a region related to the two big clouds. The best matching localization was the space between them. **Figure 13** shows the measurements during the second experiment. After 40 min the amplitudes were of a size as found in any natural spread-F. In conclusion, a gross ionospheric instability arising from the non-linear development of a bubble did not occur, but equatorial spread-F was for the first time artificially stimulated.

There was another outcome of the Colored Bubbles campaign. It was derived from the observed motions of the five small releases on upleg, which exhibited a strong vertical variation of the horizontal plasma flows. In an attempt to understand the reasons for the shear flows, an analytical model using flux tube integrated quantities was developed and the temporal evolution of the post-sunset equatorial ionosphere computed. Thereby it was possible to analyze the roles of the three most important contributors to shear flow and post-sunset rise of the F-region,

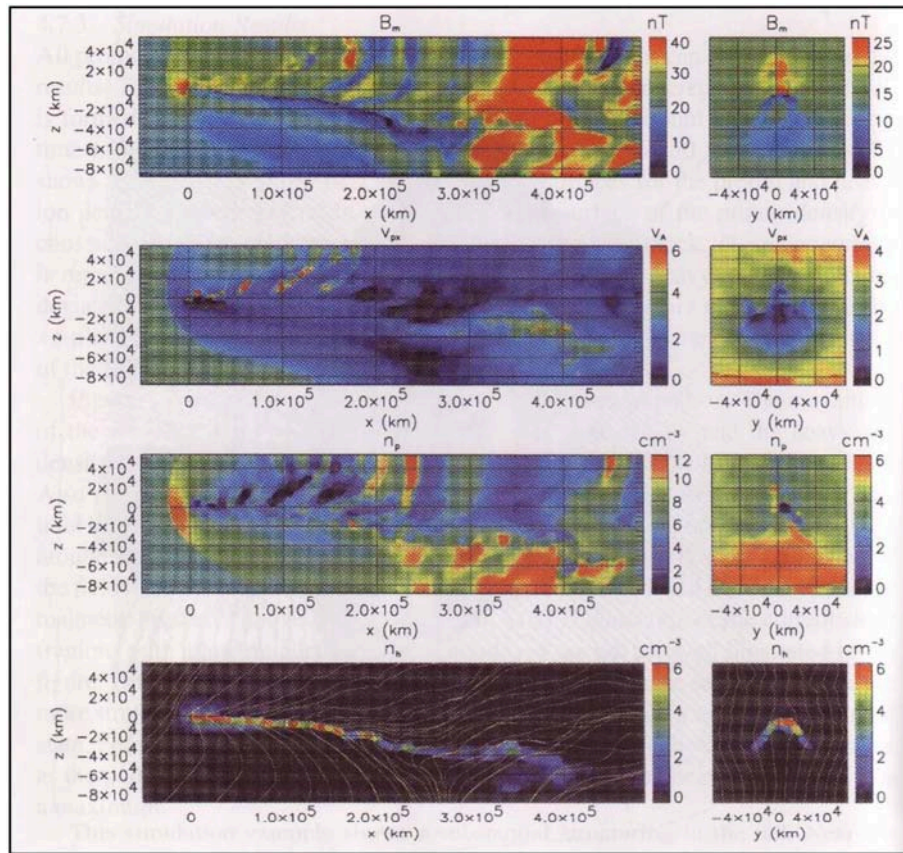


FIGURE 12 | 3D bi-ion fluid simulation box at $t = 200 \Omega_h^{-1}$ of magnetic field, B_m , proton velocity, v_{px} , proton density, n_p , and heavy ion density, n_h , for a 5,000 km large mass loading event in the solar wind. Left: xz-cut, right: yz-cut. Field lines in the bottom xz-cut are projections of the 3D field through the tail (Christian Fischer in Szegő et al., 2000).

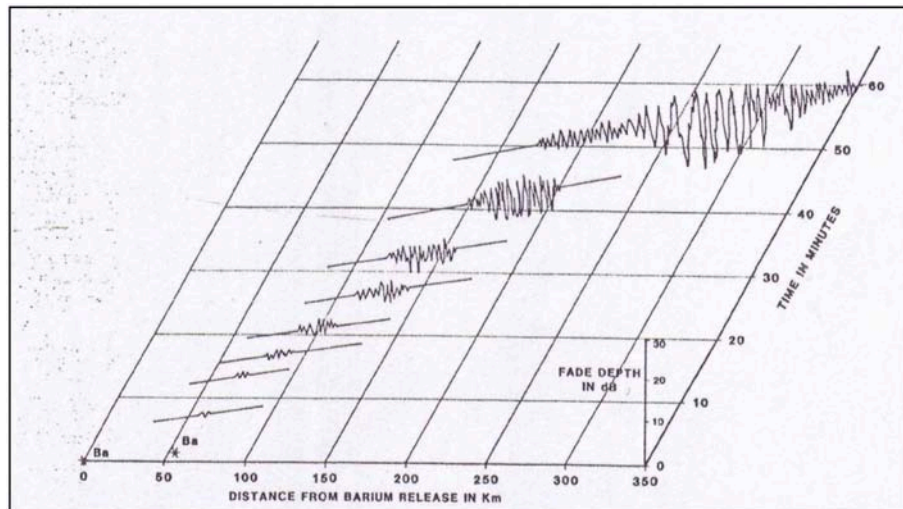


FIGURE 13 | Scintillation amplitudes of a radio beacon signal through a region disturbed by two injected barium clouds at 350 km height propagating eastward (Reproduced from Johnson and Hocutt, 1984 with permission of John Wiley and Sons).

the neutral wind dynamo, the Hall current, and the divergence of the equatorial electrojet during sunset (Haerendel et al., 1992).

The author is not aware of any subsequent experiment with the same goal and employing plasma clouds. Attempts with generating ionospheric holes are being reviewed by Paul Bernhardt (this volume).

Magnetospheric Modifications

Subsequent to the classical paper by Kennel and Petschek (1966) on the electromagnetic interaction of energetic particles trapped in the Earth's magnetic field with ion cyclotron and whistler waves, a great many of publications appeared dealing with the possibility of modifying artificially the resonant energies by artificially changing the ambient density. The first suggestion came from Neil Brice (1970). Many detailed calculations followed showing that plasma enhancement could lower the resonant energies to an extent that the originally stable ions or electrons were destabilized leading to enhanced precipitations and measurable auroral light emissions (Cornwall, 1974; Cupperman and Landau, 1974). When the Max Planck Institute for extraterrestrial Physics was offered a free ride on the second Ariane test launch, project "Firewheel" was initiated. A mother spacecraft carrying 12 release containers and four instrumented sub-satellites were built in labs in the USA, Canada, the United Kingdom, and Germany under the leadership of MPE. Besides the goals described above (section Interaction With the Solar Wind) in the context of the later performed artificial comet experiments, a central objective was to realize a substantial plasma density enhancements by a lithium release at about 9.5 R_E geocentric distance near midnight and observe *in-situ* the response of the wave activity and from ground the stimulated precipitation through enhanced light emissions. In the Li experiment, a volume of 500 km cross-section would receive a density enhancement of more than one order of magnitude, and the phase velocities of the cyclotron waves would be lowered by the same ratio. Not discussed was the possibility of a maser effect by trapping of the waves. To the great disappointment, not only of the participants but also of a widely interested plasma physics community, the Ariane L02 launch failed.

While the AMPTE mission did not allow for releases inside the magnetosphere, there appeared a new opportunity with the Combined Release and Radiation Effects Satellite of NASA which was launched in July 1990. Besides a diagnostic payload the spacecraft carried 24 chemical release containers for 14 different experiments. A larger number of research institutions and scientists joined the observational efforts and divided the responsibilities for each of the 6 types of investigations. Two releases of lithium at high altitudes conjugate to the auroral ionospheres were planned and performed. However, no positive results have been reported.

Proof that the predicted stimulated precipitation can work in principle is supported by a natural experiment, the so-called SubAuroral Morning Proton Spots (SAMPS) (Frey et al., 2004). Following a geomagnetic storm, the previously strongly eroded plasmasphere usually starts growing in size because of refilling with cool ionospheric plasma. Still existing fluxes of energetic protons from the ring current then drift into a

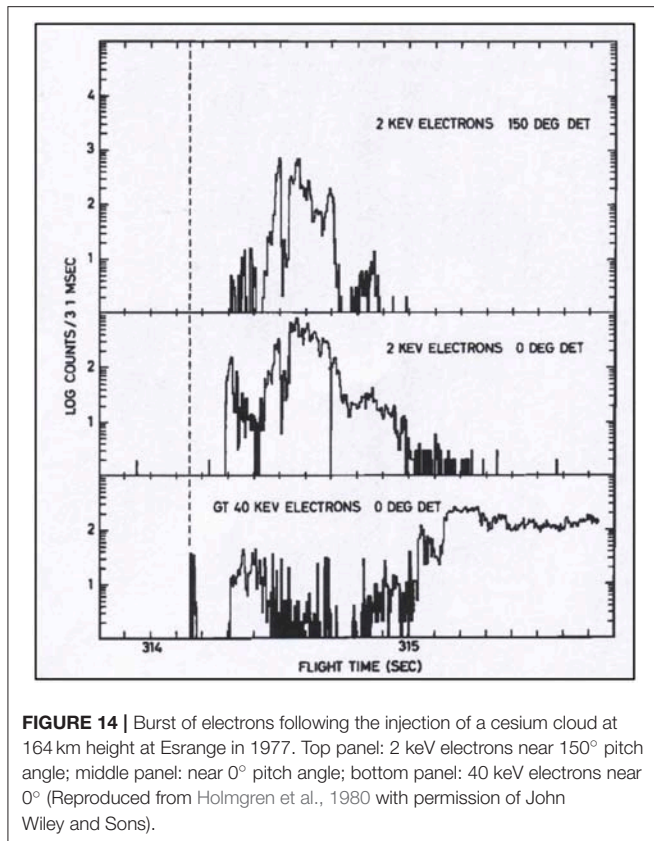
new high-density environment lowering the threshold for ion-cyclotron resonances. This was the situation reported in the cited paper. With no diagnostics available testifying for the stimulated wave-particle interactions, only the appearance of isolate proton emissions in the morning sector had to be taken as evidence.

Auroral Stimulation Experiments

By contrast with the unsuccessful attempts to stimulate auroral particle precipitation in the magnetosphere, there was a rich and successful activity to do so in the auroral ionosphere. Probably the first encounter of electron precipitation stimulated by a plasma injections occurred during two barium experiments in Esrange/Kiruna. Observing auroral emissions from the E-region, Stoffregen (1970) found enhancements of the 5,577 and 4,278 Å lines by a few percent lasting for a few tens of sec. No other diagnostics being available this was only indirect evidence for stimulated electron precipitation.

Kelley et al. (1974) reported about perturbations connected with the barium injection process performed on five rocket flights in the auroral ionosphere dedicated to comparing different techniques of measuring electric fields. Seven types of perturbations were identified including changes in electron temperature and density, in the ac and dc electric fields, and in the electron flux. Triggering these effects had not been planned for. The configurations cloud-payload were therefore somewhat accidental. While the perturbations on the local plasma environment are to be expected, changes in the energetic electron fluxes could be indications of long-range effects. However, in the light of later experiments (see below) it was rather puzzling that in the only once observed enhancement of the electron flux, lasting for about 20 s, the pitch-angle distribution above 15 keV was shifted toward 90°. The authors suggested that Alfvén waves triggered by the explosion might have interacted resonantly with trapped electrons several 100 km above the rocket.

The first experiment expressly dedicated to triggering electron precipitations was performed by Holmgren et al. (1980). It was quite appropriately named "Trigger" and consisted of a mother-daughter payload. The cesium release from the daughter on the downleg occurred 1.5 km above the instrumented mother payload. Cesium was chosen in order to obtain a fast ion production from the explosion, which for obvious reasons was planned to be performed in darkness. Forty milliseconds after the explosion a strong electric signal appeared, reaching a peak of 200 mV/m after 140 ms, and subsiding 100 ms later. About coincident with the electric field maximum a sudden increase of electrons fluxes by several orders of magnitude was registered by the downgoing as well as by the upgoing channels, with no increases in the 90° pitch-angles (**Figure 14**). Most pronounced were the fluxes at 2 keV. The spectra obtained by Lundin and Holmgren (1980) are most revealing. The sharp fall-off above a few keV, together with the field-aligned concentration, means that there must have been electrostatic acceleration along $\underline{B}$. Downgoing fluxes of up to 11 erg cm⁻² s⁻¹ corresponded to a current density of 1 μA/m².



Marklund et al. (1988) analyzed the Trigger results and concluded that the expanding neutral cloud swept up the ambient plasma creating a radial Hall current which was partially balanced by a reverse Pedersen current driven by the generated polarization field. The net current was closed by field-aligned currents in both directions. Thus, two Alfvén wave were launched. Estimating the Pedersen and wave conductances the authors concluded that the expanding neutral cloud, including a preceding shock wave, constituted a voltage generator. On the basis of the measured amplitude of the electric field and the dimensions of the cloud, field-aligned current densities well above those observed in auroral arcs can be derived, although with a high degree of uncertainty. Some of the authors therefore discussed the possibility of the current becoming unstable and developing anomalous resistivity somewhere in lower densities above the cloud (Kelley et al., 1980; Yau and Whalen, 1988). This way field-parallel voltages of 2 kV and above could have been generated. However, the above cited energy flux would be inconsistent with such a model and also with the presence of similar electron fluxes from below. Dissipation of the field-aligned current due to the finite parallel resistivity of the ionosphere and runaway electrons emerging therefrom, have not been considered. On the other hand, it is very unlikely that the field-parallel electron fluxes could have been generated by pitch-angle scattering of magnetospheric electrons. This hypothesis was therefore discarded by Holmgren et al. (1980). Equally difficult to explain is that next to the

dominant electron component of 2 keV short enhancements were measured as well at energies up and above 40 keV and near 0° within a second after injection. The experiment left many open questions.

Next to the electron enhancements and the electric field pulses also electric field and density waves were observed, first overlapping with the electron burst and, after 1 s, for another interval of about 1 s (Kintner et al., 1980). The second pulse in the range of 1–2 kHz appeared when the Cs-cloud had surrounded the payload. The authors interpreted the waves as excited by an ion-ion streaming instability between Cs^+ and the ambient ions.

Years later the Swedish group tried to repeat the Trigger experiment with two rocket flights, the Tor project, however with less success (Holmgren et al., 1988). With similar intentions Wescott et al. (1985) injected barium with shaped charges transverse to the magnetic field in sunlight for photo-ionization. Electric field pulses as well as waves were observed but no stimulated electron precipitation. However, promptly following a barium release, ions up to 6.8 keV with auroral intensity were observed, not energized, but presumably pitch-angle scattered by the explosion.

Another type of active experiments along the same lines were the plasma depletions caused by the exhaust of the Space Shuttle (Mendillo et al., 1975) and, intentionally, by releases of H_2O , CO_2 , and N_2 from explosive charges of nitromethane and ammonium nitrate in Project Waterhole (Yau et al., 1981; Whalen et al., 1985). The goal was to reduce the local plasma density by a factor of 10 and, in the presence of strong field-aligned currents, lower the threshold for current instabilities. The ensuing effects were to be observed *in situ* and by watching the auroral emissions along the affected flux tube. A release into a stable auroral arc led to large depressions in the local plasma density, in the precipitating electron flux, and in the emissions at 5,577 Å. This lasted for more than 2 min. In another rocket flight, the electron flux at 1.5 keV decreased slightly, but large enhancement were observed at small pitch angles between 0.1 and several keV. Furthermore, large transient electric and magnetic fields appeared with the passage of the shock front from the explosive release. The strong transverse perturbation field of 275 nT clearly indicated the generation of field-aligned currents. Yau and Whalen (1988) speculated that electrostatic ion cyclotron waves may have been excited by the current in the reduced plasma density of the hole.

The cited paper by Yau and Whalen (1988) contains an excellent review of all the here mentioned auroral modification experiments including the relevant associated literature. What stands out among the various attempts, is the Trigger experiment of 1977. The mechanism leading to enhanced electron fluxes, most likely accelerated out of the ionospheric plasma background, remains largely unexplained. At the time of writing of this review, active plasma experiments in space appear to be stories from the pioneering past. However, the open questions remain und would deserve new attempts for their clarification. However, one must acknowledge that, like in case of the magnetospheric modification experiments, it is not easy to initiate nature to give away its secrets by the modest means accessible to space researchers.

SPECIAL PHYSICAL PROCESSES

Critical Ionization Velocity Experiments

Alfvén (1954) introduced the critical ionization velocity (CIV) effect in a theory of planetary formation. It is a beam-plasma discharge developing when a beam of neutral particles crosses a magnetized plasma. The energy is drawn from the kinetic energy of the neutrals transverse to the magnetic field in the plasma frame. The transverse velocity component must exceed a critical threshold, v_{crit} :

$$v_{n\perp} > v_{crit} = \sqrt{\frac{2eV_{ion}}{m_n}} \quad (9)$$

(m_n is the mass of the neutral particle). The idea is that instabilities of the freshly generated ions excite plasma waves (e.g., lower hybrid), by which the electrons are resonantly heated up to and beyond the ionization potential, V_{ion} , of the neutrals. If the electrons make at least one ionizing collision during the time of contact with the neutrals the discharge can proceed. The latter is the so-called *Townsend condition*. It implies a lower limit on the density of the neutrals.

Alfvén's idea soon found much interest among theorists and experimenters, foremost in Alfvén's laboratory in Stockholm. Whereas, theories concentrated on the energy coupling processes from ions to electrons, various situations in astrophysics, cometary and planetary sciences were proposed to be sites of this effect. It was soon recognized that space offered good possibilities for testing the effect in the absence of walls and with the availability of *in-situ* diagnostics of particles and waves without disturbing the process. It so happened that in the 1970's a series of rocket flights was to be performed in Esrange/Kiruna (Project Porcupine). The payloads, highly instrumented for plasma research in the auroral ionosphere, also carried a shaped charge experiment (see section Diagnostics of Electric Fields by Tracing Visible Plasma Clouds). The author, PI of the project, decided to dedicate one of the shaped charge firings to studying the CIV effect (Haerendel, 1982).

Figure 15A explains the situation. The neutral barium jet contained a large fraction of atoms with velocities well above v_{crit} for barium, which is 2.7 km/s. In order to separate CIV ionization from photo-ionization, the injection occurred below the terminator for the lowest ionizing photons from the sun. The charge was pointed at an angle of 28° with respect to B , so that, after a short while, initially ionized barium would appear in sunlight and could be well-separated from subsequently photo-ionized particles. **Figure 15B** shows the result of the experiment of 1979 in terms of densitometer traces of the cloud as observed from ground: A narrow streak along the field lines through the injection point and a diffuse cloud to the right. Haerendel (1982) estimated that 30 % of the component above v_{crit} , i.e., between 15 and 20% of the total, had been ionized initially. Since the main payload was offset from the flux tube through the Ba jet, only fringe effects of the generated electromagnetic perturbations could be observed. The presentation of the observations was accompanied by

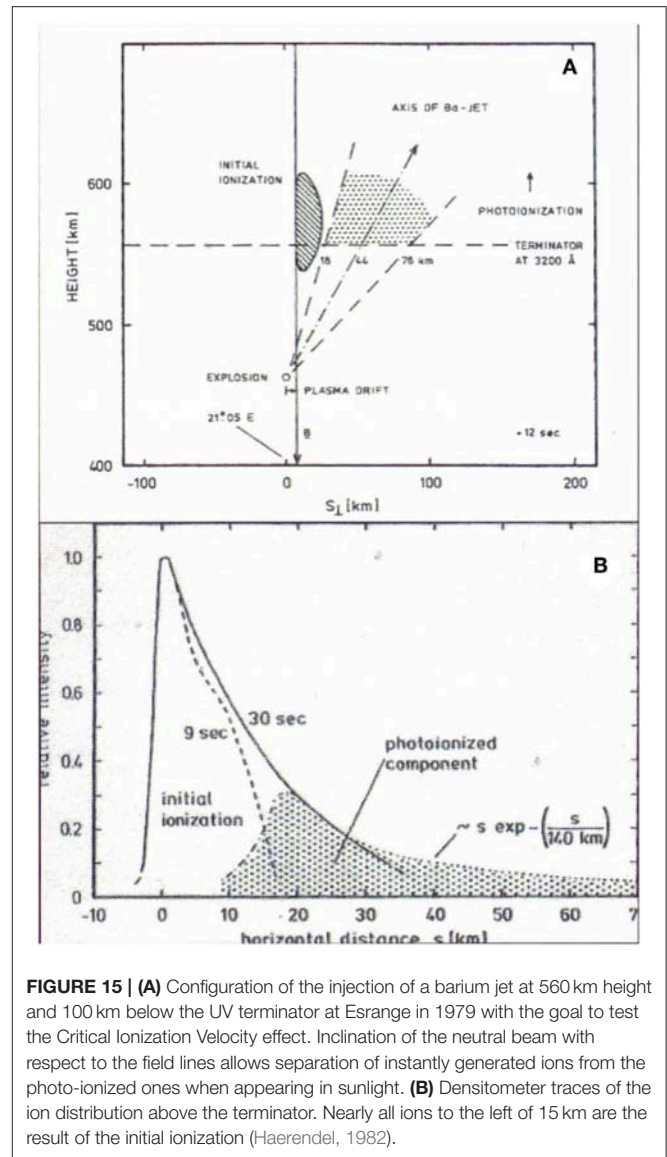


FIGURE 15 | (A) Configuration of the injection of a barium jet at 560 km height and 100 km below the UV terminator at Esrange in 1979 with the goal to test the Critical Ionization Velocity effect. Inclination of the neutral beam with respect to the field lines allows separation of instantly generated ions from the photo-ionized ones when appearing in sunlight. **(B)** Densitometer traces of the ion distribution above the terminator. Nearly all ions to the left of 15 km are the result of the initial ionization (Haerendel, 1982).

theoretical considerations of the energy transfer process and of the limitations on the successful burning of the discharge. Both aspects were subsequently taken up in a great number of similar space experiments and theoretical investigations. Excellent reviews have been published by Brenning (1992), Lai et al. (1996), and Lai (2001).

The disturbing finding was that none of the following rocket experiments up to CRIT 2 in 1989 showed any CIV associated yield. Several theories tried to explain the reason for the discrepancy with the Porcupine result (e.g., Machida and Goertz, 1988; Torbert, 1988; Papadopoulos, 1992; Moghaddam-Taaheri and Goertz, 1993). However, finally the CRIT 2 rocket experiment (Torbert et al., 1992) and several releases from the CRRES satellite (Wescott et al., 1994) produced yields of at least a few percent. In CRIT 2, for instance, the release-rocket configuration was such that the energized electrons could be intercepted by the main payload directly along the field line of

the explosion and compared with the optical yield. Still many questions remained with respect to the enormous differences found for the CIV related ionization between the various experiments. The cited references and many more to be found in these papers demonstrate to how much depth the exploration of the CIV process in space, laboratory, and natural environments has been subjected to carefully designed experiments, theory, and numerical modeling.

“Skidding” of Fast Plasma Clouds

The Combined Release and Radiation Effects Satellite (CRRES) was designed to perform a number of plasma injection experiments with diverse objectives, most of them along the lines of previous studies, such as diamagnetic cavities (cf. section Magnetic Cavities), magnetosphere-ionosphere coupling (cf. section Magnetosphere-Ionosphere Coupling), field line tracing and acceleration (cf. section Field-Parallel Acceleration), stimulating wave-particle interactions in the magnetosphere (cf. section Magnetospheric Modifications), and testing the CIV effect (cf. section Critical Ionization Velocity Experiments). However, a new goal was conceived under the name “skidding.” It was to study the initially delayed braking of a plasma cloud injected from an orbiting spacecraft into the upper ionosphere (350–450 km). What is new in comparison with the coupling of ion clouds to the ionosphere as discussed in section Relative Motions of Plasma Clouds and Auroral Arcs, is the impact of the high inertia of the injected ions. Mitchell et al. (1985) had studied the situation with numerical simulations and predicted that a 10 kg cloud would be able to skid for tens of seconds and tens of kilometers before coming to rest with respect to the ionosphere. The ion braking sets up a polarization field with $\vec{E} \times \vec{B}$ in the direction of the neutral cloud motion. Other than in the magnetosphere-ionosphere coupling discussed in section Magnetosphere-Ionosphere Coupling, the ambient ionosphere not only transfers momentum. Propagating almost instantly to the ionosphere, momentum is to be shared with the mass along the flux tube and with the region of current closure in the lower ionosphere. The motion decays due to dissipation by the Pedersen current, but with a delay because of the inertia temporarily stored in the background. The CRRES results (Huba et al., 1992) confirmed the predictions, however, skidding times and distances turned out to be lower than calculated, namely 3 s and 9 km for a 10 kg release. Further discussions on the current state of skidding theory and experiment can be found in a review paper by Winske et al. (2019).

Cycloid Bunching

At one of the CRRES barium releases at an altitude of 6,180 km a new phenomenon was observed, the formation of non-field-aligned curved structures (Bernhardt et al., 1993). The effect is explained as consequence of the motion of ions generated from a fast moving neutral cloud expanding as a spherical shell. If the linear motion of the spacecraft is faster than the cloud’s expansion velocity, the trajectories of the injected ions are dominated by cycloidal motions transverse to $\vec{B}$, while also proceeding along the field. At the cusps of the cycloids the transverse velocity goes to zero and the density is strongly enhanced. This leads

to phase bunching as a purely kinetic effect, no instability involved. However, there may be modifications caused by the resulting electric polarization fields. The key issue underlying the appearance of visible structures is the Doppler shift and the deep absorption profile of the solar Fraunhofer line at 4,555 Å. Depending on the component of the gyromotion along the cloud-sun line, the brightness is varying with the cyclotron frequency. Unfortunately, the observing intensified CCD camera did not have the sensitivity to resolve the variations with gyro-frequency, however 2 sec exposures showed for about 10 s curved structures, alluding to the contour of the expanding neutral cloud. Bernhardt et al. (1993) supported the interpretation as due to cycloid bunching by an integral solution of the Vlasov equation for the evolution of the ion distributions including the Doppler effect. The observed phenomenon once more demonstrated that with space plasmas single particle effects can be important, in this case for the optical appearance, in the artificial comet experiment even for the overall dynamics.

CONCLUSIONS

Letting the here reported physical processes passing one’s mind, one is induced to ask three questions: What has been and will remain unique? What has been superseded by later investigations with diagnostic instrumentation? And what is worth doing again? It is undisputable that what was described in section Diagnostics of Electric Fields by Tracing Visible Plasma Clouds, the role of the *plasma clouds as tracers* of natural motions, has been replaced by the development of double Langmuir probes for measuring electric fields. By contrast, the detailed insights into *coupling processes in magnetosphere and solar wind* summarized in section Coupling to the Ionosphere and Interaction With the Solar Wind are hard to obtain from orbiting vehicles in combination with ground-based measurements. *Modifications of the ionosphere* with plasma injections have only been partially successful. In the magnetosphere they did not succeed at all. Would it make sense to try again? The efforts would have to be substantially higher than those made in the past. On the other hand, the equatorial ionosphere and, in particular, the spread-F instabilities, have been thoroughly investigated by radar and *in-situ* rocket and satellite measurements. The physics of *wave-particle interactions in the magnetosphere* has been and still is being explored by many satellite missions in close cooperation with theory and simulations. In both cases there is little reason for new active experiments. The situation differs with respect to the *stimulation of electron precipitation* in the auroral ionosphere. Only few incomplete data sets are available. Little is really understood. There is no other way to clarify what is happening than with active experiments. By contrast, the *CIV experiments* with rockets and satellites have generated plenty of significant data and have stimulated intensive support by lab experiments and theory. Further space experiments are not likely to add much new information. *Skidding and cycloid bunching* are processes of minor general importance and the few experiments done can be considered as sufficiently well-understood.

We passed over the *coupling processes* treated in sections Coupling to the Ionosphere and Interaction With the Solar Wind without conclusions. The experiments were certainly unique and generated a wealth of new insights. It is hard to imagine how they could be replaced by diagnostics from orbiting satellites. Plasma cloud experiments supported by *in situ* and remote observations seem to be the only way. Can we say that there is enough material available and sufficiently developed understanding? While magnetospheric and ionospheric experiments have been continued into the early nineties by the CRRES mission, artificial comet experiments were done only twice by AMPTE. Would it be a good idea to undertake a new mission? It would most likely add new facets to the solar wind-plasma cloud interactions, clarify some details, but would probably not produce new fundamental insights. It has to be taken into account that the experimental and logistic efforts are great. Raising support for a new mission from a space agency would not be easy today. On the other hand, related physical processes are being and will continue to be explored with cometary and planetary missions.

We must accept that active plasma experiments in space were pioneering deeds of the past. They led to many unexpected discoveries and explored territory unknown at that time. Those actively involved experienced a most fascinating and often adventurous period of their lives. This short review tries to direct attention to the main achievements in understanding the encountered physical processes. It cannot convey the fascination.

AUTHOR CONTRIBUTIONS

The author confirms being the sole contributor of this work and has approved it for publication.

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Active Experiments in Space: The Future

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Planned active space experiments and ideas for future active space experiments are reviewed. Three active experiments being readied are DSX (Demonstration and Space eXperiments), SMART (Space Measurement of Rocket-released Turbulence), and BeamPIE (Beam Plasma Interaction Experiment). Ideas for future experiments include relativistic-electron-beam experiments for magnetic-field-line tracing, relativistic-electron-beam experiments to probe the middle atmosphere, plasma-wave launching using superparamagnetic-nanoparticle amplification of magnetic fields, the heavy-ion mass loading of collisionless magnetic-field-line reconnection, the use of electrostatically charged tethers to pitch-angle scatter radiation-belt particles, cold plasma releases to modify magnetospheric plasma physics, and neutral-gas releases to enhance neutral-particle imaging of the magnetosphere. Technologies that are being developed to enable future space active experiments are reviewed: this includes the development of compact relativistic accelerators, superparamagnetic particle amplified antennae, CubeSats, and a new understanding of how to control dynamic spacecraft charging. New capabilities to use laboratory facilities to design space active experiments as well as new computer-simulation capabilities to design and understand space active experiments are reviewed.

Keywords: active space experiments, plasma physics, magnetospheres, ionosphere, laboratory astrophysics, space physics

1. INTRODUCTION

Space active experiments are experiments that deliberately perturb the space environment in ways that can yield new information about the environment. They offer unique ways to gather scientific information, to study the interaction between space platforms and the space environment, and to perform space engineering. Active experiments can be used to study ionospheric physics, magnetosphere-ionosphere coupling, cometary physics, and magnetospheric plasma waves. Importantly, some experiments can only be performed in space. Space-based plasma-physics and plasma-astrophysics experiments can uniquely address the physics of large-scale plasmas, long-range coupling, and truly collisionless physical processes. In general, particle distribution functions can be obtained with more accuracy and less perturbation in space experiments than in laboratory plasma experiments. Besides scientific exploration, active experiments also support national security. For instance, a motivation of future space engineering comes in the design of active experiments for radiation-belt remediation, whereby an enhanced radiation belt environment is rapidly weakened by means of an external forcing. For scientific and engineering experiments in space, there will be needs for other space experiments to gain understanding of the interaction of those scientific and space-engineering platforms with the space environment.

There has been a rich history of active experiments in space (c.f. Grandal, 1982; Winckler, 1992; Raitt, 1995; Unan and Rietveld, 1995; James et al., 1998; Haerendel, 2018; Pongratz, 2018; Prech et al., 2018; Mishin, 2019; Winske et al., 2019 for reviews). These past experiments have involved electron and ion beams, plasma releases, chemical releases, tethers, antennae, and nuclear detonations. They span several decades, starting from high-altitude nuclear detonations in the late fifties to the plasma and chemical release experiments of the mid-nineties. In more recent years, the active-experiments program has changed, focusing on ground-based modification of the ionosphere by intense electromagnetic waves from facilities like HAARP (High frequency Active Auroral Research Program) and Arecibo.

At the “Active Experiments in Space: Past, Present, and Future” workshop in September 2017 in Santa Fe, New Mexico (Delzanno and Borovsky, 2018), several planned and proposed space active-experiment missions were discussed: these and other future missions are described in sections 2 and 3 of this report (Note that, in this paper, we only focus only on space-based active experiments: We do not review ground-based ionospheric modification experiments, but we acknowledge that these experiments are and will remain a very important component of the overall active-experiments program.). Among the advantages that future space active experiments will have over past active experiments are (1) better diagnostics, (2) newer technologies, and (3) better planning via modern computer simulations. These aspects are discussed in section 4, while conclusions are drawn in section 5. Following the “mandate” from the Santa Fe workshop, the goal of this paper is to demonstrate the importance and uniqueness of space active experiments and to generate increased enthusiasm toward an area that, fostered by many new innovations, can tremendously improve our understanding of the near-Earth environment.

At the Santa Fe workshop, there was also an overwhelming call to pass the knowledge and capabilities of active space experiments on from the older generation to newer scientists.

2. PLANNED EXPERIMENTS

Three interesting active experiments (DSX, BeamPIE and SMART) are planned in the next few years and their objectives are briefly reviewed here. Note that all three experiments have a common objective to investigate wave-generation processes in space and this fits into the broader picture of how artificially-injected electromagnetic waves could be used for radiation-belt remediation (e.g., Inan et al., 2003; Dupont, 2004) or for communication.

2.1. The DSX Dipole Antenna

The Demonstration and Science eXperiments (DSX) of the Air Force (Scherbarth et al., 2009) is currently scheduled for launch by the summer of 2019 aboard the Space-X Falcon Heavy. With an orbit of 6000×12000 km, 42 degrees inclination, it will explore the Medium Earth Orbit (MEO) environment and particularly the slot region of the electron radiation belts. DSX carries an 80-m long dipole antenna, which will be the largest, unmanned, self-supporting structure ever deployed in space, and

a comprehensive suite of space environment sensors. Its primary science objective is to study Very Low Frequency (VLF) wave transmission in MEO, including the injected VLF power by antennae in space and the interaction of VLF waves with the local particles of the environment. In this regard, DSX will work in conjunction with the VLF and Particle Mapper (VPM) nanosat mission in Low Earth Orbit (LEO), which will act as a far-field probe for DSX. Conjunctions with other spacecraft and ground stations will also be pursued. The secondary science objectives are (1) to map the local MEO radiation and plasma environment and (2) collect data to understand environmental effects and the degradation of selected spacecraft electronics and materials.

2.2. The Beam-PIE Cerenkov Wave Emission

The Beam Plasma Interaction Experiment (Beam-PIE) is a suborbital rocket experiment funded by NASA and led by Los Alamos National Laboratory. Its launch is planned for the spring/summer of 2020 from Poker Flat, Alaska. Beam-PIE is a mother-daughter system (see **Figure 1**), where the mother rocket will carry a new, compact electron accelerator technology driven by high-electron-mobility transistors. The accelerator is pulsed, designed to provide tens mA of current and energies up to 54 keV. The daughter system hosts a wave receiver and particle instrument to characterize the local environment, at a distance of 1–5 km from the mother rocket. The primary objectives of BeamPIE are two. The first is to demonstrate and increase the technology readiness level of the new electron accelerator technology for space applications. The second is to study wave generation from pulsed electron beams and quantify the generation efficiency of whistler waves relative to extraordinary-mode type waves. If waves of sufficient amplitude can be generated, a secondary science objective will be the investigation of wave-particle interaction physics and the changes to the local particle populations, possibly induced by the beam-generated waves.

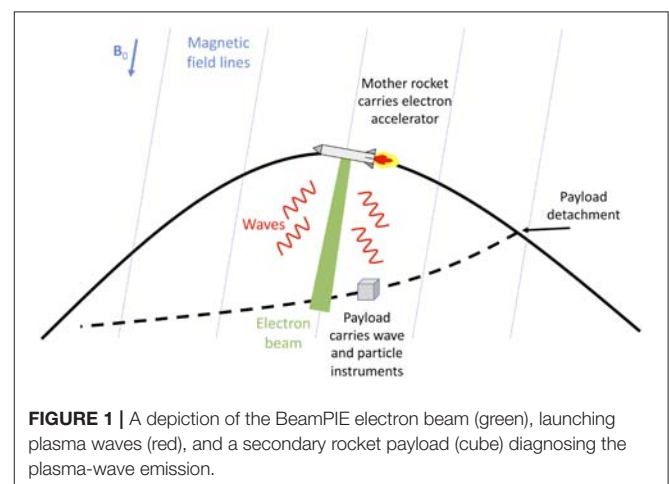


FIGURE 1 | A depiction of the BeamPIE electron beam (green), launching plasma waves (red), and a secondary rocket payload (cube) diagnosing the plasma-wave emission.

2.3. The SMART Barium Shape Charge Experiment

The Space Measurement of Rocket-released Turbulence (SMART) is a sounding-rocket experiment concept developed by the Naval Research Laboratory (Ganguli et al., 2015). At an altitude of ~ 700 km, a shaped-charge explosion will release 1.5 kg of barium atoms at high velocity (~ 10 km/s) across the Earth's magnetic field. In ~ 30 s, the barium atoms photo-ionize and create an ion ring distribution in velocity space that is unstable to electrostatic lower-hybrid waves and develops broadband lower-hybrid turbulence. SMART targets a regime of parameters where the linear damping rates are smaller than the non-linear scattering rates, implying that lower-hybrid waves can be converted into whistler or magnetosonic waves (and secondary lower hybrid waves), before significant dissipation and local plasma heating occurs. Furthermore, the electromagnetic whistler waves can propagate out of the ionospheric source region into the magnetosphere and never return to it. Estimates of the net energy extracted from the initial ring distribution (~ 5 – 10%) translate into whistler wave amplitudes of the order of 200 pT (Ganguli et al., 2015), which are easily detectable from magnetospheric spacecraft. The SMART rocket will carry the barium release module and an instrumented payload that will characterize the local turbulent source region. Operating in conjunction with magnetospheric spacecraft like THEMIS (Time History of Events and Macroscale Interactions during Substorms) to detect the SMART-induced waves, the SMART science objective is to unravel the physics of lower-hybrid turbulence in magnetized plasmas. An estimated launch date for SMART is the middle of 2021 (G. Ganguli, 2019, private communication).

3. POTENTIAL FUTURE EXPERIMENTS

At the “Active Experiments in Space: Past, Present, and Future” workshop in Santa Fe (Delzanno and Borovsky, 2018), several concepts for future space active experiments were presented, and during audience-participation discussions, the attendees highlighted the need to design active experiments to investigate (1) magnetic-reconnection onset, (2) the triggering of substorms by active experiments, (3) the mass loading of ongoing collisionless reconnection, (4) critical-ionization-velocity physics, (5) Alfvén-wave transits from one hemisphere to the other, (6) conjugate traveling-ionospheric-disturbance phenomena, and (7) magnetosphere-ionosphere coupling. There were also discussions of the pros and cons of repeating previous active space experiments with newer experimental designs and with more-powerful modern diagnostics. Calls were made by the attendees for active space experiments to address issues beyond plasma physics and the space environment: the need for experiments addressing problems in planetary physics, astrophysics, and extreme environments were suggested. Some of that workshop discussion has been incorporated into subsections 3.1–3.7 into section 4. Some of these are experiments that address large-scale issues of magnetospheric

physics, such as magnetosphere-ionosphere connectivity, triggering atmospheric discharges, triggering substorms, and producing pitch-angle scattering of magnetospheric particles into the atmosphere.

3.1. Electron Beams and Magnetic-Field-Line Tracing

The goal of this project is to accurately connect magnetospheric spacecraft measurements to ionospheric phenomena. Much of the connection between ionospheric physical processes and magnetospheric physical processes is not known. This is particularly true for the aurora and the magnetospheric processes that cause the aurora (Swift, 1978; Borovsky, 1993; Haerendel, 2011). Without understanding which physical processes act in the magnetosphere, one cannot assess the impact of auroral occurrence on the dynamics of the magnetosphere. The magnetospheric processes are unknown because the space-physics community has not been able to unambiguously connect spacecraft measurements in the magnetosphere to specific auroral forms. Magnetic field models can be used to connect large-scale regions of the magnetosphere to large-scale regions of the ionosphere (e.g., Feldstein and Galperin, 1985; Elphinstone et al., 1991; Galperin and Feldstein, 1996) but the magnetic-field models fail for the detailed mapping that is needed for auroral physics (Weiss et al., 1997; Ober et al., 2000; Shevchenko et al., 2010; Nishimura et al., 2011). The holy grail of auroral research is the low-latitude auroral arc, where one school of thought has the arcs in the ionosphere magnetically mapping out into the dipolar region of the magnetosphere (McIlwain, 1975; Meng et al., 1979; Mauk and Meng, 1991; Pulkkinen et al., 1991; Lu et al., 2000; Motoba et al., 2015), while another school has them mapping into the stretched magneto tail (Birn et al., 2004, 2012; Sergeev et al., 2012a; Hsieh and Otto, 2014). One active experiment methodology, proposed to overcome the problem of connecting magnetosphere measurements with ionospheric phenomena, is the use of an electron accelerator on a spacecraft making measurements in the magnetosphere (Borovsky et al., 1998; Delzanno et al., 2016). This is depicted in **Figure 2**. Firing the electron beam into the atmospheric loss cone and optically imaging the atmospheric beam spot using ground-based cameras can unambiguously connect critical magnetospheric measurements of plasma, flows, fields, and waves to the various auroral forms. (This spacecraft-deployed electron beam is called out in the NRC Decadal Survey (National Research Council, 2012) as a needed emerging technology for space physics.) 1 kW of beam power into the upper atmosphere will produce 3 W of optical emission in the 3914-Å band of N_2^+ (Dalgarno et al., 1965; Marshall et al., 2014). To get 1 kW of beam power, 25 mA of beam current at 40 keV is needed or 1 mA of beam current at 1 MeV is needed; firing the beam for 1 s would remove 0.025 C or 0.001 C of negative charge from the spacecraft, respectively. Spacecraft charging in the tenuous collisionless magnetospheric plasma is a potential problem. The development of compact, efficient relativistic-electron accelerators (cf. section 4.1) greatly reduces the spacecraft-charging problem by reducing the beam current.

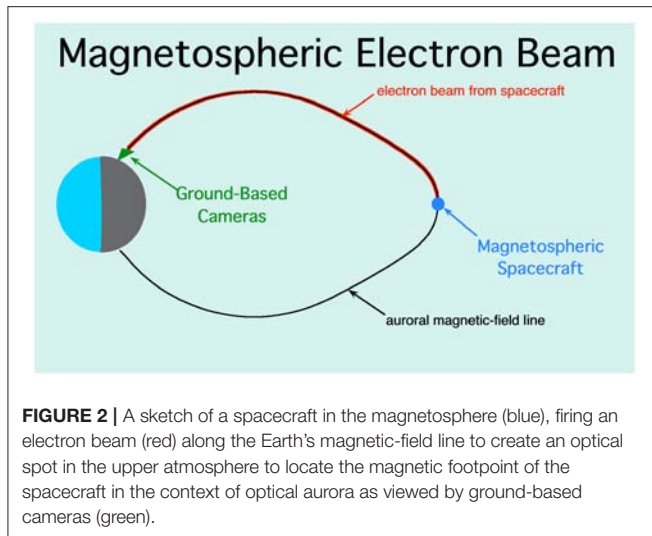


FIGURE 2 | A sketch of a spacecraft in the magnetosphere (blue), firing an electron beam (red) along the Earth's magnetic-field line to create an optical spot in the upper atmosphere to locate the magnetic footprint of the spacecraft in the context of optical aurora as viewed by ground-based cameras (green).

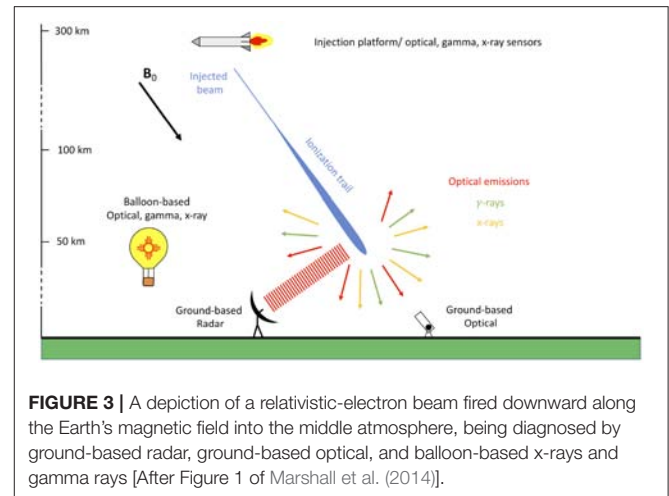


FIGURE 3 | A depiction of a relativistic-electron beam fired downward along the Earth's magnetic field into the middle atmosphere, being diagnosed by ground-based radar, ground-based optical, and balloon-based x-rays and gamma rays [After Figure 1 of Marshall et al. (2014)].

Using a plasma contactor (e.g., Olsen, 1985; Comfort et al., 1998) on the spacecraft, simulation analysis (Delzanno et al., 2015a,b; Lucco Castello et al., 2018) finds that the mechanism of ion emission from the surface of a kilometer-sized plasma-contactor plume will be able to balance the 1-mA electron-beam current and keep spacecraft charging to a low level. Magnetic-field measurements onboard the spacecraft are used to point the accelerator beam into the atmospheric loss cone. Increasing the beam energy could further reduce the beam current, which would further reduce the risk of spacecraft charging. However, for beams with energies above 1 MeV, beam pointing becomes a challenging issue, since the atmospheric loss cone shifts away from the 0° -pitch-angle direction owing to finite-gyroradius effects (Mozer, 1966; Il'ina et al., 1993; Porazik et al., 2014). The present design for the 1-MeV compact accelerator (Lewellen et al., 2019) yields an electron beam with an angular divergence of $<0.05^\circ$, including the beam's electrostatic expansion after exiting the accelerator (The space charge of the 1-MeV 1-mA beam is very low.). Such a beam easily fits inside an atmospheric loss cone that is $>1^\circ$. One unfortunate fact is that the electron beam produces optical emission in the exact same airglow wavelength bands as does the electron aurora, making it difficult for the ground-based cameras to identify the spacecraft beam spot in the presence of active aurora: using a time-coded on-and-off beam sequence and looking for the blinking beam spot greatly improves the detection. There is also the possibility of detecting the beam spot via ground-based radar (Izhovkina et al., 1980; Uspensky et al., 1980; Zhulin et al., 1980; Marshall et al., 2014, 2018) and of using the relativistic beam to do ionospheric and atmospheric experiments diagnosed by the radar.

3.2. Relativistic Electron Beams Into the Middle Atmosphere

Ionospheric and atmospheric experiments could be performed with a relativistic-electron beam fired downward from the magnetosphere, or fired from a low-altitude spacecraft, a rocket (e.g., Nunz, 1990; O'Shea et al., 1991), or even from a balloon if the beam energy is high enough (See depiction

in Figure 3). Electrons with energies of a few MeV range out at about 40–50 km altitude (Marshall et al., 2014), where the atmospheric number density and collision density is about the same as in a 1-Torr vacuum chamber. Ionization and recombination/attachment experiments have been suggested by Banks et al. (1988, 1990), Neubert et al. (1996), and by Neubert and Gilchrist (2004); these experiments could be diagnosed by ground-based radar (cf. Figure 3). Issues that could be investigated include the decay of electrical conductivity, electron-attachment rates, and the transport of negative and positive ions in the atmospheric electric field (Borovsky, 2017). The stimulation of atmospheric-electricity discharges by the electrical-conductivity paths, provided by relativistic-electron-beam ionization columns above thunderstorms, have been suggested by Banks et al. (1988, 1990), Neubert et al. (1990), Neubert and Gilchrist (2004), and Marshall et al. (2018) with the discharge current flowing between the top of a thunderstorm and the ionosphere. The energy deposition of a 1 kW beam is about 50 times the energy deposition of a naturally occurring relativistic-electron microburst (Lorentzen et al., 2001; Borovsky, 2017). These triggered thunderstorm discharges could be diagnosed by ground-based optics or by ground-based electric (Thomas et al., 2004; Sonnenfeld and Hager, 2013), magnetic (Whitley et al., 2011), or electromagnetic (Rhodes et al., 1994; Qin et al., 2012) measurements. The observation of upward accelerated energetic particles from the triggered discharges (e.g., Lehtinen et al., 2000, 2001) has also been suggested by Neubert and Gilchrist (2004); such observations can be made from the spacecraft or rocket that carries the relativistic-electron accelerator. Atmospheric chemistry modification by relativistic electron beams has also been explored (Neubert et al., 1990; Marshall et al., 2018), with the suggestion of diagnostic via ground-based spectroscopy; the chemistry of NO_x , HO_x , and ozone production in the middle atmosphere by energetic electron precipitation is of particular interest for the information it can supply about the interaction of the Earth's radiation belt with the Earth's climate system (Rodger et al., 2010; Andersson et al., 2012; Verronen et al., 2013).

3.3. Modifying Magnetic Reconnection With Heavy Ions

Gaining an understanding of the factors that control the onset of collisionless reconnection and the factors that control reconnection rates is of great importance to magnetospheric physics and solar-coronal physics. Using an artificial plasma cloud to modify collisionless reconnection (to initiate the onset of reconnection or to mass load and reduce ongoing reconnection) is a possibility. The onset of collisionless reconnection is an outstanding science issue that would improve the prediction of substorm occurrence (McPherron et al., 1973; Sergeev et al., 2012b) and of solar-flares occurrence (Priest, 1986; Li et al., 2017). It has been speculated both that the introduction of heavy ions to a plasma will make it (a) easier for the plasma to reach conditions for the onset of field-line reconnection (Baker et al., 1982, 1985, 1989) or (b) harder for it to reach the onset of reconnection (Liu et al., 2013; Liang et al., 2016). The onset of reconnection in collisionless plasmas is usually thought to be caused by the thinning of a current sheet to a thickness below ion-inertial-length or ion-gyroradius spatial scales (Hesse and Birn, 2000; Liu et al., 2014). It has been variously speculated that introducing heavy ions (1) alters tearing modes that thin the current sheet, or (2) changes the ratio of current sheet thickness to gyroradii, or (3) mass loads current sheets. More simulation work with modern kinetic simulation codes (e.g., Karimabadi et al., 2011; Pritchett, 2013; Birn and Hesse, 2014) is needed to verify these conjectures. The mass loading of ongoing reconnection is an important concept for the reduction of solar-wind/magnetosphere coupling via magnetospheric feedback (Borovsky et al., 2013; Walsh et al., 2014). A cloud consisting of 1 kg of barium ions (e.g., Bryant et al., 1985), with a diameter of 1,000 km, has a mass density of about 1,000 AMU/cm³, which is about 20 times higher than the mass density of the magnetosheath plasma at the dayside reconnection site. This barium mass density is sufficient to effectively turn off dayside reconnection within the cloud, if the barium cloud could be released close enough to the dayside reconnection site. **Figure 4** depicts the fact that getting the cloud (#1, purple) over the reconnection diffusion region (red) is helped by the fact that the barium ions will be carried into the reconnection line by the Mach-0.1 inflow of ambient plasma into the line. If the barium ions at the dayside magnetosphere could be optically imaged, the reconnection rate could be gauged by the speed of the barium ions carried in the reconnection outflow. Targeting reconnection away from the nose of the magnetosphere may allow ground-based imaging of the barium cloud via cameras located beyond the solar terminator. Since the location of the reconnection X-line may be difficult to predict, experiments on the mass loading of the reconnection outflow fan (which can extend across the entire dayside magnetopause) with barium releases, may be easier to implement. This is depicted as cloud #2 in **Figure 4**. Getting barium into the reconnection fan is again aided by the Mach-0.1 inflow of ambient plasma into the fan. Comparison of Earth's reconnection regimes (with and without heavy ions) with reconnection observations by MAVEN at Mars with O⁺ and O²⁺ ions (e.g., Harada et al., 2015) and by

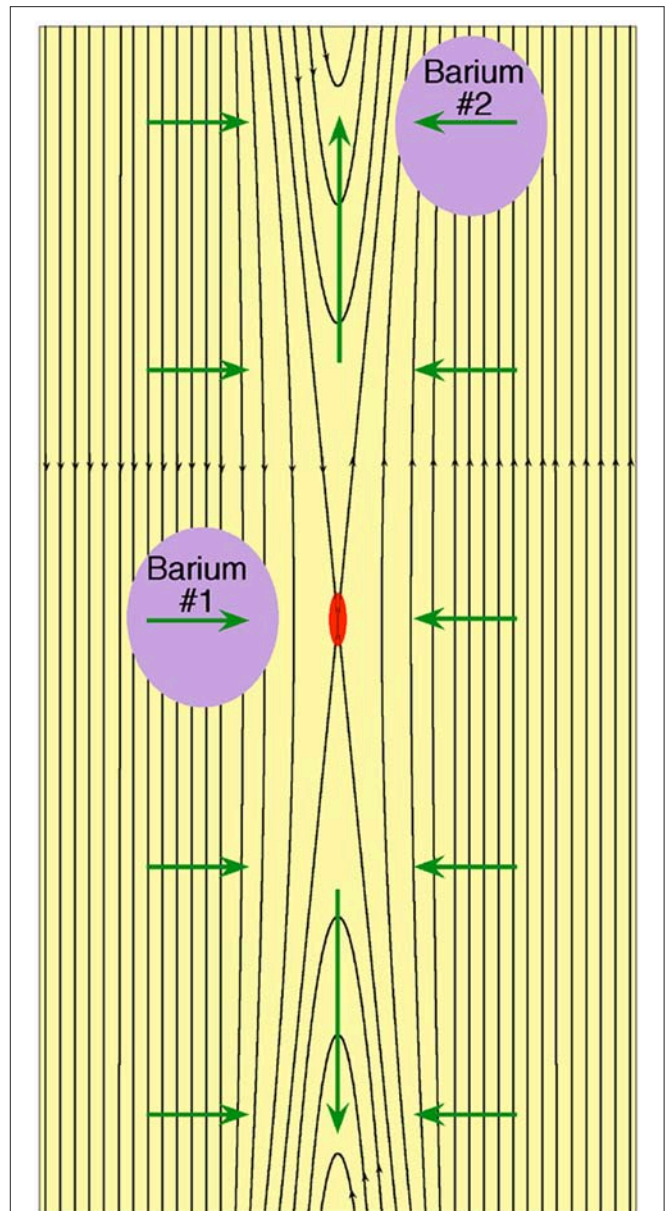


FIGURE 4 | Magnetic-field-line reconnection on the dayside of the Earth is depicted with the solar-wind plasma to the left and the magnetospheric plasma to the right. The reconnection diffusion region is marked in red. There is a flow (green arrows) everywhere into the vertical plane of the current sheet that feeds plasma into the reconnection site or into the reconnection-outflow fan. Two barium clouds are depicted: Cloud #1 is being drawn into the reconnection diffusion region and cloud #2 is being drawn into the reconnection outflow fan.

Juno at Jupiter with S⁺ ions could be useful for preparing and planning heavy-ion active experiments as described above.

3.4. Plasma-Wave Launching With Rotating-Magnet Antenna

Efficient ways to launch plasma waves into the magnetosphere are of interest for future technologies, such as radiation-belt remediation (Inan et al., 2003; Dupont, 2004). A space

experiment has been suggested (Dennis Papadopoulos, private communication 2018) for the launching of whistler, EMIC (electromagnetic ion-cyclotron), and Alfvén waves, from a low-Earth-orbit spacecraft or a rocket using a superparamagnetic-nanoparticle-amplified rotating magnetic antenna. A rotating magnetic field can be created with an orthogonal pair of magnetic coils driven by sinusoidal currents with a 90° phase difference between the two coils. At the center of the orthogonal-coil pair, a vacuum vessel containing ~1 kg of superparamagnetic nanoparticles (Raikher et al., 2004) would act to amplify the strength of the rotating magnetic field by a factor of about 100, greatly amplifying the efficiency of the coils to launch whistler waves, EMIC waves, or Alfvén waves, depending on the frequency applied to the coils. Alfvén waves are important for understanding magnetosphere-ionosphere coupling (Goertz and Boswell, 1979) and whistler and EMIC waves are important for coupling the evolution of the radiation belt to the evolution of other magnetospheric plasmas (Borovsky and Valdivia, 2018). Without the superparamagnetic nanoparticles, the two-coil rotating-magnetic-field concept has been successfully tested in the laboratory for the launching of Alfvén waves (Gigliotti et al., 2009; Karavaev et al., 2011) and whistler waves (Karavaev et al., 2010). As discussed in section 4, this proposed active experiment is being enabled by the technology development of superparamagnetic nanoparticles. A similar active space experiment has been suggested by Karavaev (2010) and de Sonria-Santacruz et al. (2014), using a mechanically rotating superconducting magnetic coil.

3.5. Space Tether Experiments

Tethers are a powerful technology tool that can be used to facilitate space experiments (Johnson et al., 2017; Huang et al., 2018): enabling multipoint measurements, launching whistler and Alfvén waves, acting as an antenna, and providing propulsion. Past active experiments using tethers (Lorenzini and Sanmartin, 2004; Cartmell and McKenzie, 2008) involved examining the dynamics and electrodynamics of tethers, the electrodynamic interaction between tethers and the space plasma environment, and the emission of plasma waves. More space experiments are needed to further understand the interactions of electrodynamic tethers with the plasma environment (e.g., Choiniere et al., 2001; Siguier et al., 2013; Janeski et al., 2015) and to explore wave launching by tethers (Estes, 1988; Lutgen and Neubauer, 1994; Sanchez-Arriaga and Sanmartin, 2010). One suggested active experiment is to use a kV-charged tether to electrostatically pitch-angle scatter radiation-belt particles into the atmospheric loss cone as the particles pass through the tether's sheath (Hoyt and Minor, 2005; Huboda de Badyn et al., 2016), although estimated time scales for remediation appear too long (~1 yr). Another interesting active experiment involving a tape tether used to explore the upper atmosphere has been suggested by Sanmartin (Sanmartin et al., 2006; Sanmartin, 2010): ambient ions would be accelerated into a long, negatively biased tape producing secondary electrons which are then accelerated off the tape to excite an artificial aurora in the upper atmosphere.

3.6. Cold-Plasma Releases

The idea of using cold-plasma releases in the magnetosphere, to trigger instabilities that stimulate electron and/or ion precipitation and produce artificial auroras, has been suggested since the seventies (Brice, 1970; Brice and Lucas, 1971; Cuperman and Landau, 1974). In the magnetosphere, EMIC waves are driven by hot-ion temperature anisotropies associated with magnetospheric convection and charge exchange, and whistler-mode chorus waves are driven by hot-electron temperature anisotropies associated with substorm injections. The addition of cold ions to the magnetosphere by a plasma release will change the growth rates and saturation amplitudes of EMIC waves (Fu et al., 2016; Gary et al., 2016). Whereas, the addition of cold electrons to the magnetosphere by a plasma release will change the growth rates and saturation amplitudes for whistler waves (Cuperman et al., 1973; Cuperman and Sternlieb, 1975; Gary et al., 2012). The cold ions and electrons also change the energetic-particle resonance conditions for EMIC waves and whistler waves, respectively (Summers et al., 1998). Provided that certain conditions on the anisotropy of the distribution function are met, a plasma injection can allow more particles to precipitate in concert with the development of the instability and the generation of electromagnetic waves. A likely location in the magnetosphere for such a cold-plasma experiment is in the nightside of the dipolar region, where there can be anisotropic hot populations to drive waves, and where ordinarily, there is an absence of cold ions and electrons owing to magnetospheric convection bringing plasma in from the magnetotail.

Magnetospheric barium and lithium release experiments were performed in the Active Magnetospheric Particle Tracer Explorer (AMPTE) (Krimigis et al., 1982) and the Combined Release and Radiation Effects Satellite (CRRES) programs, with several scientific goals including substorm triggering and stimulation of particle precipitation. In particular, three lithium releases (G-5, G-6 and G-7), by CRESS at ~33,000 km, did not show enhanced aurora that would be a sign of enhanced wave-particle interactions (Bernhardt, 1992). Two barium releases (G-8 and G-10) showed increased auroral activity within 5 min from the release, although the definitive association with the release was uncertain (Bernhardt, 1992). Similarly, a magnetotail barium release by AMPTE, during the development of a substorm, showed the barium cloud moving antisunward and was interpreted with the formation of a reconnection plasmoid (Baker et al., 1989).

Given the importance of substorms and wave-particle-interaction physics for magnetospheric dynamics, cold-plasma release active experiments should be pursued in the future with modern technology to test relevant theories of magnetosphere-ionosphere coupling.

3.7. Hydrogen-Gas Releases for Enhancing Energetic-Neutral-Atom Imaging of the Magnetosphere

Information (densities and temperatures) about the global distribution of hot plasma in the Earth's magnetosphere is

obtained by imaging the energetic neutral atoms that are produced when energetic plasma ions charge exchange with the Earth's neutral-hydrogen geocorona (e.g., Roelof et al., 1985; Gruntman, 1997). One difficulty with the neutral-atom-imaging technique is that the measured fluxes of neutral atoms are line-of-sight integrated through the entire magnetosphere. Scime and Keesee (2018) propose a method to focus the neutral-atom measurements on a single point in space by releasing neutral hydrogen gas at that point in space to greatly enhance the number of charge exchange collisions, and hence greatly enhance the flux of energetic neutral atoms originating from the release site. This would provide higher spatial resolution measurements of the magnetospheric hot plasmas of the magnetosphere at the same time as global images are being obtained.

4. CRITICAL TECHNICAL ADVANCES

For the future of space active experiments, several technical advances are being made that will facilitate new and improved experiments. Further, there is presently improved laboratory and computer simulation support capabilities for the design of future space experiments.

4.1. Advances in Electron Accelerators

For future electron-beam experiments in the magnetosphere, the research and development advances of compact relativistic-electron accelerators has been crucial. Accelerators that have relatively high efficiency (bus power to beam power) are in development (Lewellen et al., 2019): this increased efficiency saves battery weight on the spacecraft and reduces battery recharging time from solar panels, the latter enabling more beam time. The critical thermal issue of heat removal from the accelerator has been reduced by the development of a method for re-tuning the frequency fed to the linear accelerator, as the accelerator changes temperature and mechanically expands. Designs for the remote operation of fault-tolerant linear accelerators are in development.

4.2. Superparamagnetic Nanoparticles

As discussed in section 4, advances in the development of superparamagnetic nanoparticles for amplifying AC magnetic fields is making the design of more-powerful space-based wave antennas possible.

4.3. CubeSats

The development and availability of low-cost CubeSats has increased access to low-Earth orbit for experiments (Bahcivan et al., 2012; Poghosyan and Golkar, 2017) and diagnostics (Blum et al., 2013; Fish et al., 2014). Active-space-experiment diagnostics with constellations of CubeSats (Glumb et al., 2016; Deng et al., 2017) is a new possibility.

4.4. Controlling Spacecraft Charging

As discussed in section 3, the advancement in our understanding of methods to ameliorate spacecraft charging in electron-beam experiments is allowing for lower-risk experiments to be designed. A significant advance has been made by

the interpretation of plasma contactors in the collisionless magnetosphere, working as ion emitters rather than electron collectors (Delzanno et al., 2015a,b; Lucco Castello et al., 2018). This work was guided by new plasma-simulation capabilities (see section 4.6).

4.5. Laboratory Support for Developing Space Experiments

Laboratory experiments are becoming increasingly important for our understanding of plasma and space physics and in support of (active or inactive) space experiments, as reinforced in a recent review (Howes, 2018) that coined the term "laboratory space physics." Often driven by similar advances in diagnostics and technology, laboratory experiments complement space experiments by allowing a more controlled environment that can be diagnosed much more extensively. On the other hand, laboratory experiments operate with plasma densities, temperatures and, more importantly, collisionality that can be very different from those of the space environment, thus allowing scaled experiments where only ratios of relevant quantities controlling the physics of interest can be kept in the same range. Laboratory experiments are ideally suited to isolate particular physics aspects of more complex problems, while their size limitation makes it difficult to explore things like long-range coupling.

In the US, there are several facilities with a history of significant contributions to space physics and the interested reader is referred to Howes (2018) and references therein for a summary. See also Koepke (2008). Here, we only focus on the connection between laboratory and active experiments and highlight relevant experiments.

The Basic Plasma Science Facility (BAPSF) at the University of California Los Angeles is a national user facility that hosts the Large Plasma Device (LAPD), a 19-m long, 75-cm diameter cylindrical plasma column (Gekelman et al., 2016). LAPD operates with typical densities of 10^{12} cm^{-3} and electron temperatures of few eV (with lower values in the afterglow plasma). The high reproducibility of the experiments, combined with extensive diagnostics, make detailed three-dimensional characterization of the plasma an important feature of LAPD. To guide the design and interpretation of planned electron-beam experiments in space, electron-beam experiments are being performed on LAPD. While earlier experiments used a low-energy (3 keV) electron beam to explore the excitation of chirped whistler waves (Van Compernelle et al., 2015; An et al., 2016), a 1-MeV linac (Jenkins et al., 2018) is being installed on LAPD. The new experiments will study relativistic-beam stability and the generation of plasma waves, with application to solar radio bursts as well as to electron-beam active experiments for radiation-belt remediation. LAPD experiments involving a laser-generated plasma and its explosive dynamics across a magnetic field are investigating processes associated with the formation of a diamagnetic cavity and collisionless shocks (Niemann et al., 2013, 2014; Winske et al., 2019), and are relevant to early nuclear detonation experiments in space.

The Space Physics Simulation Chamber at the Naval Research Laboratory, shown in **Figure 5**, also allows for studies across different parameter regimes targeting ionospheric and magnetospheric conditions. Examples include the role of shear-driven ion-cyclotron waves in ion heating and initiation of ionospheric outflows (Amatucci et al., 1998), electron-ion hybrid instabilities important for the plasma sheet boundary layer (Amatucci et al., 2003), and the generation of electromagnetic ion cyclotron waves through shear flows (Tejero et al., 2011). More recent experiments have focused on non-linear scattering processes, successfully demonstrating the conversion of electrostatic lower-hybrid waves to electromagnetic whistler waves above an amplitude threshold (Tejero et al., 2015). This is a key aspect of the non-linear weak-turbulence physics that the SMART barium-release experiment aims to demonstrate (cf. section 2.3).

The 6m × 9m Large Vacuum Test Facility (LVTF) and the 2m × 0.6m Cathode Test Facility (CTF) at the University of Michigan's Plasmadynamics and Electric Propulsion Laboratory (PEPL) (Gallimore et al., 1996; Gilchrist et al., 2002) have been used for experimental validation of spacecraft charging mitigation induced by high-power electron beams. LVTF is capable of reaching 10^{-7} (10^{-8}) Torr and is the biggest vacuum chamber in the US. In the LVTF experiments, an isolated hollow-cathode represents the spacecraft. The hollow cathode emits a high-density charge-neutral plasma (known as the plasma contactor), while the emission of the spacecraft electron beam is mimicked through a separate power supply operated in constant-current mode. Several Langmuir probes, emissive probes and a retarding potential analyzer provide measurements of key quantities, identified by the space-experiment modeling work (Delzanno et al., 2015a,b; Lucco Castello et al., 2018). Remarkable agreement between theory and experiments has been obtained (Miars et al., 2018), thus validating the ion-emission model for spacecraft-charging mitigation for the operation of electron-beam experiments in the low-density magnetosphere (cf. section 4.4).

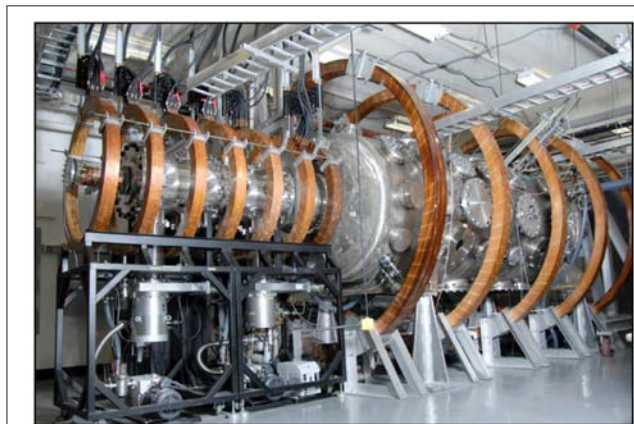


FIGURE 5 | The 7.6-m long Space Physics Simulation Chamber at the Naval Research Laboratory in Washington DC (Photo courtesy of Erik Tejero).

A Community-Coordinated Modeling-Challenge Facility that uses laboratory facilities at West Virginia University, combined with high-performance-computing modeling from interested parties, is also being proposed to study spacecraft-environment interactions (Koepke and Marchand, 2017).

4.6. Simulation Support for Designing Space Experiments

Another major advance in support of the design and planning of (active or inactive) space experiments comes from numerical simulations. This is the result of both the increased power and availability of modern high-performance computers, and also of the recent advances in development of new numerical algorithms to tackle the multiscale nature of plasmas. The major challenge comes from the large spatial and temporal scale separation typical of magnetized, collisionless plasmas. This occurs already at the microscopic/kinetic level, due to the mass difference between electrons and ions, but quickly becomes overwhelming when one compares microscopic scales with system scales.

Recent advances in the development of kinetic Vlasov-Maxwell solvers include the implicit particle-in-cell (PIC) method (where implicit refers to the temporal discretization of the method) (Chen et al., 2011; Markidis and Lapenta, 2011) and the use of discontinuous-Galerkin discretization techniques (Juno et al., 2018). Moreover, electrostatic PIC methods that employ some form of non-uniform mesh (either conforming or through adaptive mesh refinement, structured, or unstructured) are commonly used to study dynamic spacecraft-environment interactions (Mandell et al., 2006; Roussel et al., 2008; Marchand, 2012; Delzanno et al., 2013; Meierbachtol et al., 2017).

In terms of global codes for large-scale dynamics, hybrid (kinetic ions and fluid electrons) codes, running on high-performance computing platforms, are now routinely applied to study the dynamic of the Earth's magnetosphere (Karimabadi et al., 2014; Lin et al., 2017; Palmroth et al., 2018). Furthermore, methods for “fluid-kinetic coupling” are also being developed for large-scale simulations that include microscopic physics. One approach is based on a regional kinetic code locally embedded in a large-scale fluid-like simulation (which is typically from a magnetohydrodynamic code) (Sugiyama and Kusano, 2007; Kolobov and Arslanbekov, 2012; Daldorff et al., 2014; Tóth et al., 2016; Ho et al., 2018). This approach has been successfully applied to study flux-transfer events and Earth's dayside reconnection (Chen et al., 2017). Other approaches are based on higher-order fluid moments with suitable closures (Wang et al., 2015). A new method that encompasses both techniques described above has been developed in the SpectralPlasmaSolver (SPS) code (Delzanno, 2015; Vencels et al., 2016). It is based on a spectral expansion of the plasma distribution function in Hermite functions, such that the low-order terms of the expansion are akin to a fluid description of the plasma, while kinetic physics is retained by adding (possibly locally in the simulation domain) more terms to the expansion. As such, fluid-kinetic coupling is an intrinsic feature of SPS, but the method is not constrained to a fixed number of moments and the transition between fluid and kinetic regimes can be handled as

smoothly as necessary. SPS has been successfully applied to the turbulent cascade in the solar wind (Roytershteyn and Delzanno, 2018; Roytershteyn et al., 2019). Global space weather models, such as the SHIELDS (Space Hazards Induced near Earth by Large Dynamic Storms) framework (Jordanova et al., 2018), are now beginning to incorporate some of these innovations (which also include data-assimilation techniques to assimilate available observational data) and will be very important in the future to put spacecraft observations into better context, particularly for geomagnetically active times.

Finally, besides some of the more technical innovations highlighted above, we mention the Community Coordinated Modeling Center (CCMC, <https://ccmc.gsfc.nasa.gov/index.php>) (Bellaire, 2006; Rastatter et al., 2012), which hosts a large number of heliospheric, magnetospheric, and ionospheric simulation codes and models, and offers free “runs on request” using the computational resources of the center. CCMC’s goal is to provide access to modern space science simulations for the international research community.

5. CONCLUDING REMARKS ON THE FUTURE

There are still many open questions that need to be answered by future active experiments. Three examples from three research fields are given to highlight the breadth of future active experiments. For plasma astrophysics: (1) Under what conditions does the critical-ionization-velocity effect operate? For space physics: (2) What is the magnetic-field connectivity between ionospheric regions and processes and magnetospheric regions and processes? For space engineering: (3) What is the most effective way to generate various types of plasma waves from a space platform? There are also technology capabilities that need to be developed via space experiments: e.g., (i) radiation-belt remediation and (ii) power transmission between Earth and space. And there are also new, modern technologies (in a broad sense that encompasses also diagnostics, laboratory experiments and computer simulations), perhaps best exemplified by the fact that a Tesla automobile is currently traveling in deep-space orbit (Chang, 2018), that justify new and more ambitious active experiments.

In addition, active experiments that are not necessarily associated with plasma or space physics will also be extremely important. An example is the Stratospheric Controlled Perturbation Experiment (SCoPEX, <https://projects.iq.harvard.edu/keutschgroup/scopex>) experiment, which plans to release

aerosols in the stratosphere as a possible way to reduce or eliminate ozone loss and mitigate global warming.

Active experiments have a rich history of important contributions to the field of space physics. As the spiral of knowledge advances, revisiting active experiments holds a key to finally closing fundamental questions.

Some of these grand-challenge problems can only be addressed successfully with a broad cross-disciplinary team at the intersection between theory, modeling, observations, experiments (in laboratory and, ultimately, in space) and, importantly, technology. It is, however, extremely hard to develop and maintain these large collaborations until suitable opportunities open up. One potential remedy and recommendation would be to reinvigorate and expand the active space-based experiments program, which flourished in the 1970s and 1980s to test basic scientific ideas and new technologies in space, but it has reduced its footprint in recent decades (Delzanno and Borovsky, 2018).

For the field of space active experiments, the future looks busy.

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JB and GD shared equally in the planning and outlining of the manuscript and in the researching and writing of the manuscript.

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Space-Borne Electron Accelerator Design

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Renewed interest in active experiments with relativistic particle beams in space has led to the development of solid-state radio-frequency (RF) linear accelerators (linac) that can deliver MeV electron beams but operate with low-voltage DC power supplies. The solid-state RF amplifiers used to drive the accelerator are known as high-electron mobility transistors (HEMTs), and at C-band (5–6 GHz) are capable of generating up to 500 watts of RF power at 10% duty factor in a small package, i.e., the size of a postage stamp. In operation, the HEMTs are powered with 50V DC as their bias voltage; they thus can tap into the spacecraft batteries or electrical bus as the primary power source. In this paper we describe the initial testing of a compact space-borne RF accelerator consisting of individual C-band cavities, each independently powered by a gallium nitride (GaN) HEMT. We show preliminary test results that demonstrate the beam acceleration in a single C-band cavity powered by a single HEMT operating at 10% duty factor. An example of active beam experiments in space that could benefit from the HEMT-powered accelerators is the proposed Magnetosphere-Ionosphere Connection (CONNEX) experiment (Dors et al., 2017).

Keywords: electron accelerators, space-borne accelerators, radio-frequency linac, high electron mobility transistors, particle beams in space, magnetosphere, ionosphere

INTRODUCTION

The interconnection between the magnetosphere and the ionosphere has been a topic of intense research for decades. However, detailed understanding of the processes responsible for a variety of aurora activities is currently lacking. For instance, we still do not have satisfactory answers to the questions: What creates the aurora? How are the auroral ionosphere and night side magnetosphere connected through its time-varying magnetic field? What magnetospheric processes and conditions produce particular auroral and ionospheric signatures? What are the ionospheric signatures of specific magnetospheric regions, boundaries, and events? The CONNEX proposal seeks to answer these questions and establish an unambiguous connection between the magnetosphere and ionosphere through an active mapping technique using relativistic electron beams with beam energy of about 1 MeV (Dors et al., 2017). Such an experiment will be the first of its kind to use high-energy, MeV electron beams as an active probe for doing space science.

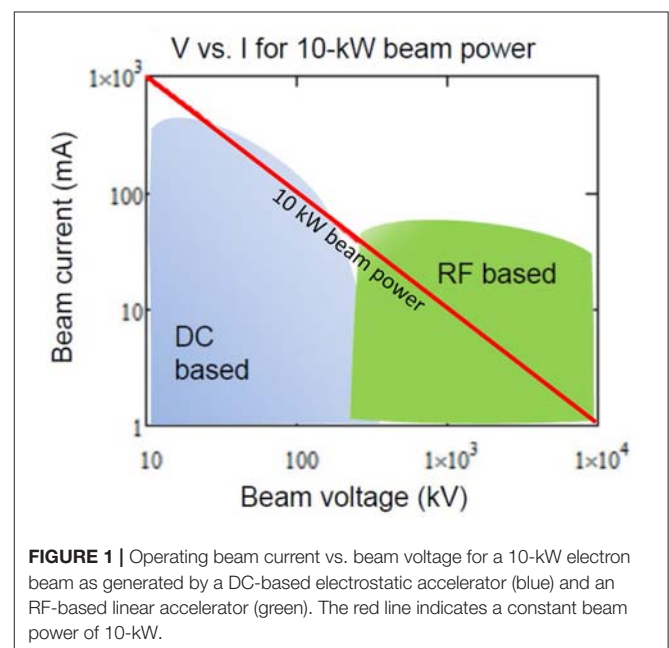
Electron beams for space experiments have previously used direct current (DC) electrostatic accelerators to deliver electron beam pulses at beam energy up to 40 keV using standard high-voltage DC power supplies. These DC electron generators are simple to design and very efficient

at converting electrical power into beam power. The first round of active beam experiments in space in the 1970s used one of these DC electron generators mounted on a sounding rocket to inject a low-energy, high-current electron beam into the ionosphere to study the interaction of electron beam with the nearby and distant magnetosphere (Hendrickson et al., 1975, 1976; Winckler et al., 1975). Subsequently in the 1980s a series of experiments were performed with an electron accelerator on Spacelab-1 with the aim of studying the interaction between the electron and plasma beams with the surrounding plasma (Obayashi et al., 1982). Many of these early experiments were performed with low-voltage, high-current electron beams in the ionosphere where the positive charge left on the spacecraft after the emission of electrons, known as spacecraft charging, was neutralized by the return current from the surrounding plasma. As the beam experiments move higher into the magnetosphere, the surrounding plasma density is reduced and charge neutralization from the surrounding plasma becomes less effective, resulting in arcing and payload failures due to severe spacecraft charging (Cohen et al., 1980; Sasaki et al., 1986). Recent efforts to mitigate the spacecraft charging problem have focused on (1) operating the accelerator at a higher beam voltage to reduce the current emitted from the spacecraft while maintaining constant electron beam power, and (2) deploying a plasma contactor to provide the surrounding plasma density necessary for the return current to neutralize the spacecraft (Lucco Castello et al., 2017). Compared to DC electrostatic accelerators, radio-frequency linear accelerators can deliver much higher beam voltage (energy) and also better beam quality, i.e., lower divergence, as well as delivering a flexible beam pulse format that can facilitate the detection of the visible light or RF signals produced by the electron beam pulses. **Figure 1** plots the range of beam current and voltage for a 10-kW electron beam using a typical DC-based electrostatic accelerator (blue) and an RF-based linac (green). The red line represents a constant 10-kW power in the electron beam. For the same beam power, the higher-voltage RF-based linac requires lower beam current, resulting in less severe spacecraft charging.

The development of RF-based particle accelerators for space missions dates back to the 1980s when Los Alamos National Laboratory successfully launched and operated a radio-frequency quadrupole (RFQ) accelerator aboard a rocket (O'Shea, 1990). The RFQ accelerated H^- ion beams that were then neutralized to produce neutral hydrogen atom beams for the BEAR (Beam Experiment Aboard a Rocket) project as part of the Neutral Particle Beam program. RF linac use time-varying electric fields along the axis of a resonant RF structure consisting of a number of RF cavities to accelerate charged particle beams. An RF cavity is a hollow piece of electrical conductor enclosing an evacuated volume that stores electromagnetic energy in the form of time-varying electric field (pointing along the cavity axis) and magnetic field (circulating near the outer cavity walls). The two halves of an RF cavity made out of copper are shown together with an HEMT (small white square) mounted on a printed circuit board in **Figure 2A**. Typical amplitudes of the accelerating electric fields range from a few megavolts per meter (MV/m) for a low-gradient structure to more than 100 MV/m for a high gradient

linac. In order to realize these high accelerating fields, RF linac have historically been driven by high-power RF sources, such as klystrons, that are capable of delivering 5–100 MW of RF power over the duration of a few microseconds. These high-power sources are large and heavy, and they require pulse forming networks and high-voltage (e.g., 50–100 kV) power supplies. The output of the source is typically shared between a large number of RF cavities, depending on the particulars of the design. For our mission to deploy a compact and lightweight accelerator in space, we need a new source of RF power that eliminates the need for high voltages and bulky pulse forming network. **Figure 2B** shows the model of a 55-cavity, 1.7-m long accelerator that weighs about 127 kg including the weights of all low-voltage RF power sources and beam control systems.

Compact RF power source now exists with the recent release of high-power solid-state RF amplifiers such as the Wolfspeed/Cree CGHV59350 high-electron-mobility transistors (HEMTs) (Cree, 2018). These HEMTs are capable of ~ 500 W of RF power each, and they can be used to power individual accelerating cavities with independent phase and amplitude controls (Lewellen et al., 2016; Nguyen et al., 2018). Our accelerator design differs from the more traditional approach of combining a large number of solid-state RF amplifiers into a high-power all-solid-state RF system (Di Giacomo, 2009). By using HEMTs for direct pumping without the power combiner and operating at relatively low accelerating gradient, we improve the efficiency of converting electrical power into electron beam power. The overall wall-plug efficiency for the HEMT-powered accelerator is estimated at 10% or greater. For the CONNEX experiment, to produce 1 kW of electron beam power (1 MeV at 1 mA average current), the accelerator is expected to have DC power consumption of 10 kW during a 10-s burst every 5 min. The average power consumption during a 4-h engagement is only about 500 W.



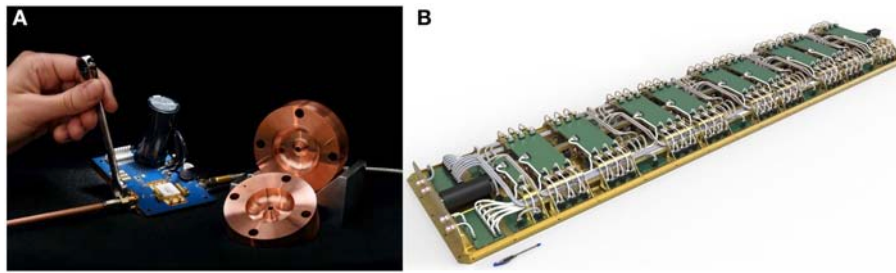


FIGURE 2 | (A) Photo of two halves of a C-band cavity and an HEMT on a printed circuit board; **(B)** A compact 55-cavity, 1.7-m long, 1-MeV electron accelerator based on HEMT solid-state RF sources.

A NEW CONCEPT FOR ELECTRON ACCELERATORS

Several factors need to be considered when selecting the accelerator technology for space applications: size, weight, power requirement (efficiency) and reliability of the accelerator system. Typical terrestrial RF linear accelerators consist of a string of either copper or niobium resonant cavities assembled in a continuous structure known as an accelerator module. Water-cooled copper accelerator modules operate at room temperature, driven by high-power klystrons or solid-state power-combined sources, and are temperature-stabilized to within a fraction of a degree C. Typical efficiency of converting electrical power to beam power is around 7% (Has Tajar et al., 2016). Niobium linear accelerator modules can approach 50% efficiency and can be driven by low-voltage, solid-state RF sources. However, the niobium cavities must be maintained below 4 K (liquid helium boiling point) to remain superconducting, and thus require a large and vibration-sensitive liquid helium cryoplant and cryomodels. Including the cryoplant power requirements, the wall-plug efficiency of superconducting accelerator modules is also usually in the single digits (Has Tajar et al., 2016). Both types of RF linac have significant size, weight, and power requirement as well as a single point of failure: the klystron and its high-voltage power supply for copper accelerators and the liquid helium cryoplant for niobium accelerators. Neither would be ideal for a space-borne accelerator.

In the past 3 years, Los Alamos National Laboratory has developed a new configuration of space-borne accelerators based on a new class of solid-state RF sources that are sufficiently small and lightweight to be deployed in space and that run on low-voltage power supplies. Our new concept of space-borne electron accelerators differs in several respects from conventional electron accelerator design, reflecting the very different environment in which it must operate. First, each accelerator cavity is individually powered by solid-state HEMTs serving as its own RF amplifier chain, and operates at relatively low gradients of 1–5 MV/m. Secondly, there is no active temperature stabilization. Instead, the cavity temperatures are allowed to rise during operation and the rates of temperature rise in individual cavities depend on the interior surface ohmic losses and the heat capacity of individual cavities. Thirdly, the cavity frequencies are monitored and adjusted with the use of active

frequency control to allow the cavities to operate over a range of temperatures.

The new accelerator configuration has several operational benefits. First, using HEMTs as the RF amplifiers running on low-voltage DC power supplies eliminates the problems of operating high-voltage devices in space. Secondly, the system wall-plug efficiency can be much higher than conventional linear accelerators, because (a) no power is expended on active cooling of the accelerator cavities, and (b) the cavities operate at relatively low accelerating gradients allowing a greater fraction of the RF power to be delivered to the beam. Finally, the accelerator system is robust against failure of individual components due to the inherent modularity of the design. For instance, in a conventional accelerator a klystron failure will definitely lead to a system shutdown, whereas in the new modular design, the failure of a single HEMT would result in only a small reduction in the total beam energy.

A key feature of the new design is the low accelerating gradient and thus a higher fraction of the RF power going into the beam. As shown in equation 1 below, operating at low accelerating gradient (E_0) and high beam current (I_b) reduces the RF power delivered to the cavity (the first term on the right-hand side of Equation (1) and increases the power delivered to the beam (the second term in Equation 1).

$$P_{RF} = \frac{|E_0|^2}{R_s} L_{cav} + V_b I_b \quad (1)$$

Here, P_{RF} is the total RF power required, R_s is the shunt impedance of the cavity per unit length, L_{cav} is the cavity length, V_b is the voltage gain of the beam through the cavity, and I_b is the beam current. For illustration, let us consider an RF accelerator design capable of generating a 1 MeV, 10-mA beam operating at a 10% duty cycle (1 kW average power) as required to effectively probe the coupling between the Earth's magnetosphere and ionosphere (Marshall et al., 2014; Dors et al., 2017). If we select an accelerating gradient of 1.5 MV/m for the C-band cavity with 1.3 cm active length—the cavity length is chosen to match the average velocity of the sub-relativistic electrons throughout most of the cavities—then the cavity power (the first term of Equation 1) is about 300 W and the voltage gain per cavity is 20 kV. With 10 mA instantaneous current, the instantaneous beam power is 200 W, so 40% of the incoming RF power (~ 500 W) is converted

into beam power. While the DC-to-RF conversion efficiency of individual HEMTs is at least 50%, since we have to use two HEMTs for each cavity due to their low gain, the DC-to-RF conversion efficiency drops to about 25% for the pair. Thus, the net efficiency of converting DC electrical power to beam power is 10%, which is still higher than a typical efficiency of terrestrial RF linacs.

A unique feature of the RF linac is its ability to produce a beam pulse format consisting of a series of pulses, minipulses and micropulses. For an RF linac operating at 5.1 GHz, the micropulses are separated by 0.196 ns, the inverse of 5.1 GHz. The length of the minipulses is set by the duration of the RF amplifier pulses, which for HEMTs is about 100 microseconds (us). During the 100-us minipulse, the beam power shall be 10 kW (1 MeV, 10 mA). Using the 25% DC-to-beam conversion efficiency, the DC power requirement for the space-borne accelerator would be 40 kW. The average power requirement would be lower since the HEMTs operate at 10% duty factor, i.e., the minipulses shall be on for 100 us and off for 900 us. The CONNEX mission requires an electron beam pulse, consisting of approximately fifty minipulses, that is sufficiently long (~ 0.5 s) to deposit a substantial amount of energy (~ 500 J) from the electron beam into the upper ionosphere to achieve good signal-to-noise ratio on the ground detectors.

The CONNEX accelerator design consists of a low-voltage DC electron gun, a buncher RF cavity, where the electron beam undergoes density modulations and forms short bunches, and 54 accelerating cavities assembled in nine groups of 6 cavities separated by focusing solenoids (**Figure 3A**). The first cavity acts as a “buncher” cavity to modulate the continuous electron beam from the DC electron gun into short bunches separated by one RF period. For most of the accelerator, the electron beam travels at sub-relativistic velocity, i.e., the particle velocity is much less than the speed of light. In a terrestrial RF linac, this would require adjusting the cavity length to match the velocity of the particle beam. In our space-borne accelerator, we shall fix the cavity to a length corresponding to the average beam velocity (about 0.4 times the speed of light). In operation, we shall adjust the RF phases of individual cavities such that the electron bunches arrive at the longitudinal center of the cavities when the accelerating field is at or near the maximum. This is made possible by using a low-level RF control system that independently phases the RF input to the HEMT amplifiers that power each individual cavity. Beam dynamics simulations using the GPT particle-pushing code (van de Meer and De Loos, 2001) show that 50% of the electrons from the DC gun are bunched, captured and accelerated continuously to 1 MeV with the use of this independent RF phase adjustment (**Figure 3B**). In addition to independent phase adjustments, the field amplitude of these cavities can also be independently controlled to maximize the capture efficiency and the total energy gain. The choice of fixed cavity length simplifies the cavity design and fabrication, and allows the heat load to remain the same for all cavities, an important feature when operating these cavities without active cooling as it simplifies the frequency stabilization.

GALLIUM NITRIDE HIGH ELECTRON MOBILITY TRANSISTORS

Wide-bandgap GaN-based HEMT are a new class of RF power devices that have recently found widespread use in wireless and satellite communication. These HEMT devices can also be used as high-power RF amplifiers over a broad range of radio-frequencies thanks to their large breakdown voltage and high electron velocity (Mishra et al., 2008). The fabrication of HEMTs typically involves growing GaN films via epitaxial layer growth on semi-insulating SiC substrates and then a thin layer of AlGaIn is grown over the GaN film to form an AlGaIn/GaN heterojunction (**Figure 4A**). Due to the different energy bandgap structures of AlGaIn and GaN, large energy band bending occurs at the heterojunction, creating a potential difference that results in a flow of free electrons (**Figure 4B**) toward the underlying GaN, forming a two-dimensional electron gas (2DEG) (Lee, 2014). This high-density accumulation of free electrons, combined with the high polarization field at the sharp interface between AlGaIn and GaN layers, is responsible for the high electron mobility in HEMTs.

HEMTs can be constructed to operate over a broad range of frequencies, with center frequencies ranging from 1.2 to 9.6 GHz and bandwidths up to $\sim 20\%$, and delivering RF power up to 700 W per device at 2–4 GHz. The RF power needed to drive a single accelerator cavity scales with the square of the accelerating field, inversely with the cavity shunt impedance per unit length (R_s) and proportionally with the cavity length, as shown in Equation 2.

$$P_{cav} = \frac{|E_0|^2}{R_s} L_{cav} \quad (2)$$

We selected the Cree HEMTs at 5–6 GHz because these HEMTs provide the highest available RF power for the electron beam, defined as the difference between the HEMT output and the cavity power. The cavity power is calculated from the expected shunt impedance for copper cavities at different frequencies assuming 1 MV/m as the accelerating gradient in these cavities. The scaling of cavity shunt impedance (a measure of how efficiently RF cavities utilize RF power in establishing the cavity accelerating field), cavity length and cavity power with frequency is shown in Equations 3–5. The HEMT output and calculated cavity power for the frequency range 2–10 GHz is plotted in **Figure 5**.

$$R_s \propto f^{\frac{1}{2}} \quad (3)$$

$$L_{cav} \propto f^{-1} \quad (4)$$

$$P_{cav} \propto f^{-\frac{3}{2}} \quad (5)$$

The RF power available for the electron beam is the difference between the HEMT output power (**Figure 5**, red dots) and the cavity power (**Figure 5**, blue curve). As can be seen in **Figure 5**, the available RF power for the particle beams is greatest at the 5–6 GHz frequency band (C-band).

To characterize the RF performance of commercial C-band HEMTs for accelerator operation, we set up a test fixture and

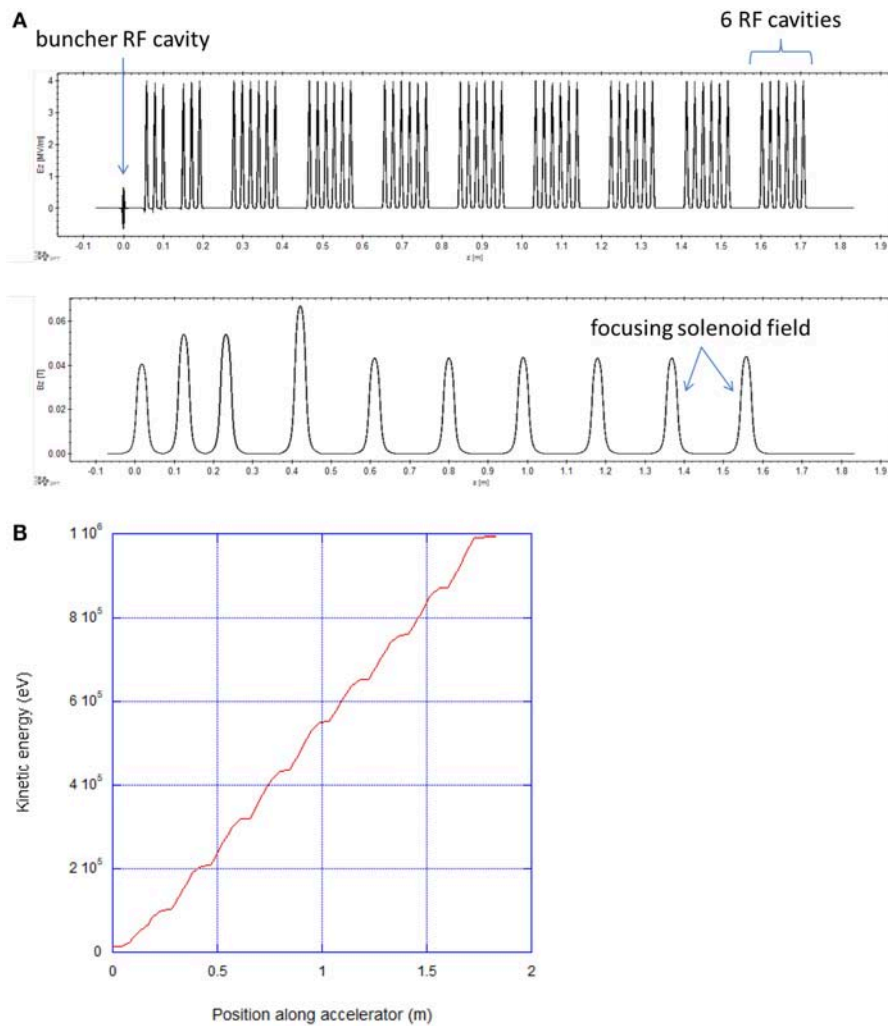


FIGURE 3 | (A) Plot of accelerating field magnitudes of individual RF cavities and focusing solenoid fields along the accelerator; **(B)** plot of kinetic energy in eV vs. position along the accelerator.

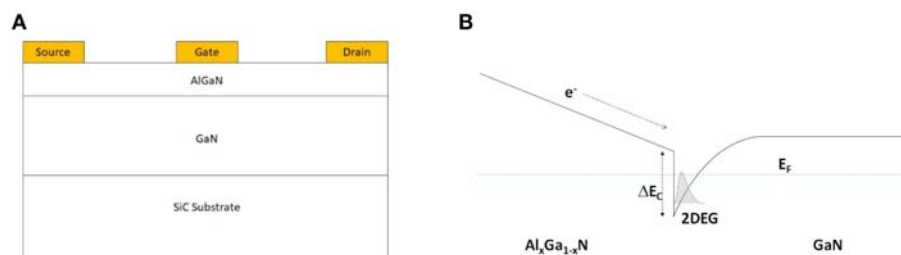


FIGURE 4 | (A) Schematic of a typical GaN HEMT cross-section showing a thin (~ 20 nm) AlGaIn layer and a micron-thick GaN layer grown on top of a SiC substrate; **(B)** The energy band structure of AlGaIn/GaN showing the formation of the two-dimensional electron gas (2DEG) at the interface.

measured the output power, small-signal gain and harmonic content of the HEMT output. The RF input with amplitude of about 0 dBm (1 mW) was generated by a low-noise continuous-wave (CW) network analyzer followed by an RF switch to

produce low-amplitude RF pulses, with duration up to 500- μ s and repetition rates up to 600 Hz, to be amplified in the preamp with 40–50 dB gain. The 40-W output from the preamp was amplified in a GaN HEMT with 10 dB small-signal gain to

produce 400 W of RF power. **Figure 6** shows the schematic of the HEMT RF accelerator performance characterization set-up. The combination of pulse duration and repetition rate allowed us to explore HEMT performance up to 30% duty factor, three times higher than the nominal rating of the device.

The output power from a single HEMT is plotted vs. the network analyzer power in dBm in **Figure 7A**, showing linear response over a range of input power until the output is saturated at about 400 W. At saturation, the single-pass gain of the HEMT is only 10 dB which requires us to have two GaN HEMTs operating in series for each cavity. **Figure 7B** shows the typical waveform of a 500- μ s RF pulse from the HEMT. The HEMT output pulse shows a power drop from $\sim$ 550 W at the leading edge to $\sim$ 400 W at the trailing edge. We have not ascertained the cause of this power drop. However, we expect the HEMT output power to depend on temperature of the AlGaIn/GaN junction; as the output power exceeds 500 W, the AlGaIn/GaN junction temperature rises. The AlGaIn/GaN HEMT drain current has been shown to drop as a function of junction temperature and one expects the output power to also decline at high temperature (Wang et al., 2013).

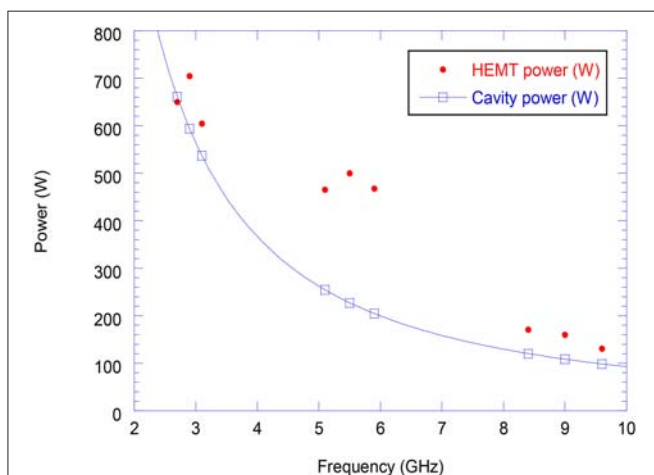


FIGURE 5 | Plot of HEMT output power (red) and calculated cavity power (blue) vs. frequency.

Several aspects of HEMT performance are of particular concern for space-borne applications. These include basic performance (power output, small signal gain, etc.), operating in low-power-consumption modes, and power drop over the pulse. The nominal requirement for a single cavity for the CONNEX project is the acceleration of a 10-mA beam through a 20-kV gap; thus 200 W of RF is needed for the electron beam power per cavity. Using Equation 2 and the measured shunt impedance of the C-band cavity, we estimate approximately 250 W of cavity power is needed to generate the required 20-keV acceleration, so each cavity will require a total of 450 W of RF power. The nominal minimum output of the HEMT is 350 W; in practice, we find HEMTs can usually produce 450 W even when operated at 50 V DC, the lower end of their operating range.

We tried to maximize the HEMT output power by adjusting the drain-source voltage (**Figure 8A**). As the drain-source voltage was increased from 40 to 90 V, the HEMT output power rose to a maximum of 610 W at 65 V and then decreased at higher drain-source voltage. We also tried to reduce the HEMT quiescent current (thus improving the average efficiency) by operating the HEMT at two sub-threshold gate bias voltages. At gate bias voltage more negative than -3.7 V, the quiescent current decreased to zero, and the HEMT power draw was zero without RF. As we increased the input RF power to 35 dBm, the HEMT generated power with 10 dB small-signal gain at -3.7 V gate bias voltage (**Figure 8B**). At -5 V bias voltage and high input power, the small-signal gain rose to more than 9 dB if the input RF power exceeded 44 dBm. These results suggest that the HEMT output power and efficiency can be improved by optimizing the drain-source and gate bias voltages.

INITIAL ACCELERATOR OPERATION AND ENERGY MEASUREMENTS

Calorimetric Measurements

We operated the single-cavity accelerator without active cooling for extended durations, measured the absorbed RF power and compared the results with RF measurements. The cavity temperatures were plotted vs. time (**Figure 9**, blue curve) and from the temperature rise, we estimated the RF power absorbed in the cavity using a calorimetric model. As the cavity resonant frequency shift is inversely proportional to the temperature rise,

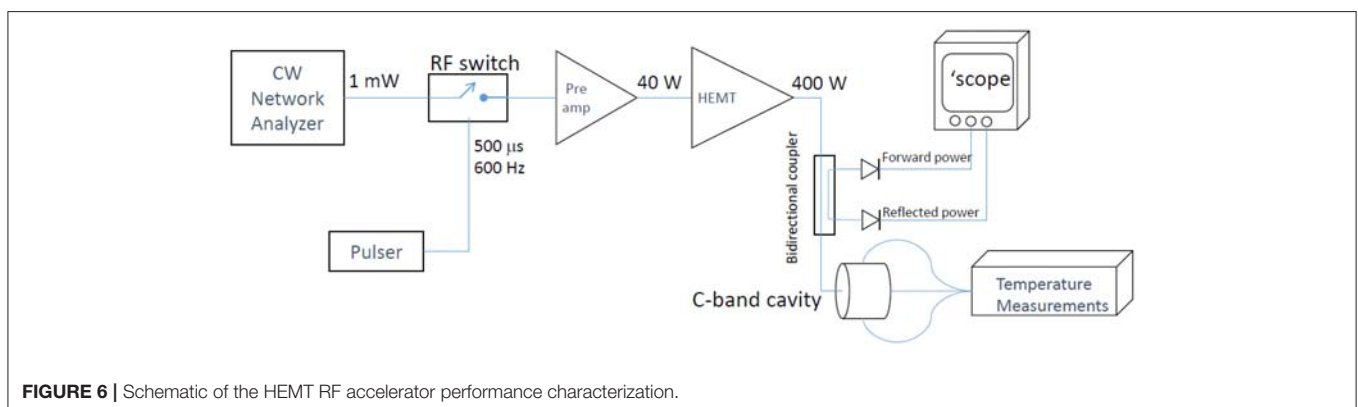


FIGURE 6 | Schematic of the HEMT RF accelerator performance characterization.

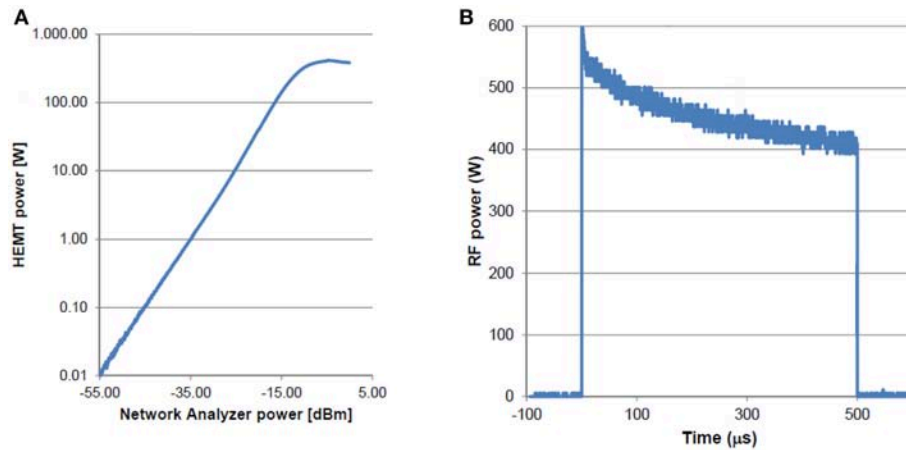


FIGURE 7 | (A) Plot of HEMT output power vs. the network analyzer power in dBm; **(B)** a typical 500-μs waveform of the HEMT output showing ~30% power drop.

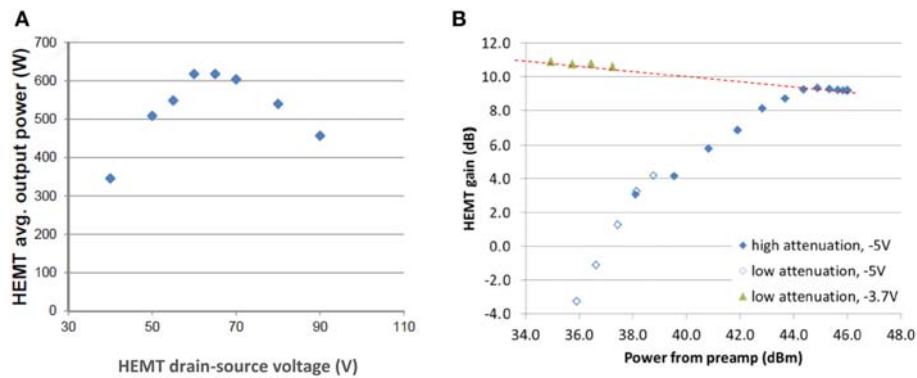


FIGURE 8 | (A) Plot of HEMT output power vs. HEMT drain-source voltage; **(B)** HEMT gain at two different gate bias voltages, -3.7V and -5V , both below the threshold to turn on the HEMT.

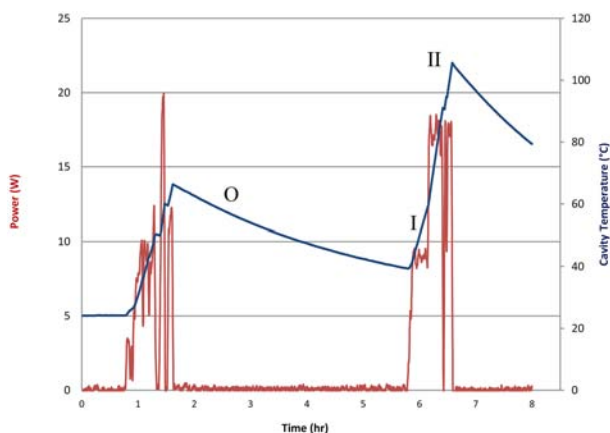


FIGURE 9 | Plots of average RF power (red) and cavity temperature (blue) as functions of time. The time period O corresponds to no RF power; I corresponds to RF on at 5% and II at 10% duty factor.

the RF source frequency was also varied to track the cavity frequency. In our simple calorimetric model, the rate of cavity heating due to the average RF power absorbed in the cavity is the sum of two terms: (1) the rate of heat causing the cavity temperature to rise, i.e., the cavity heat capacity term, and (2) the heat loss due to conduction to the surroundings, which is expressed as the inverse of the thermal resistance. The average RF power absorbed in the cavity is given by Equation 6,

$$P_{RF} = mC_v \dot{\theta} + \frac{\theta}{R_T} \quad (6)$$

where θ is the difference between the current and initial temperatures, $\theta = T - T_{ini}$, m is the copper cavity mass, C_v is copper heat capacity and R_T is the thermal resistance. The thermal resistance is calculated from the thermal decay time constant τ , defined as $\tau = mC_v R_T$, which can be extracted from the temperature decay curve (**Figure 9**) after the RF power was turned off.

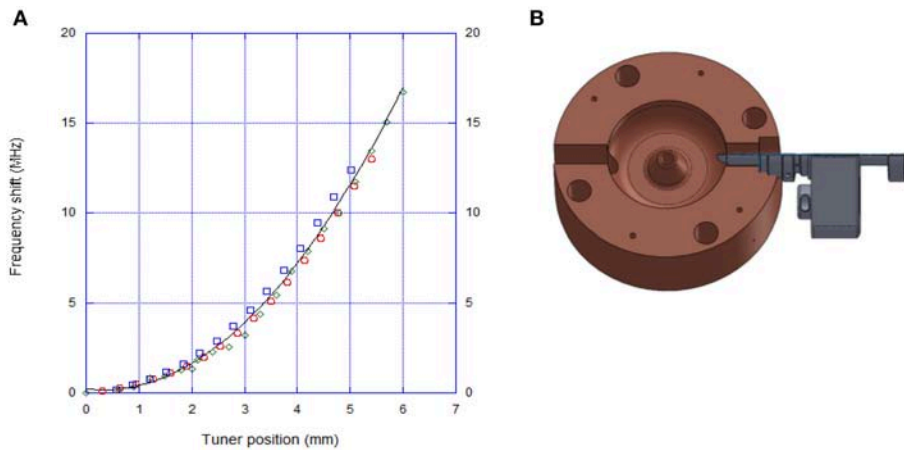


FIGURE 10 | (A) Cavity frequency shift vs. tuner position for two separate sweeps of the tuner (blue and green markers) and CST simulation results (red markers) superimposed with the fit of all three data; **(B)** Solid model of the cavity with a section view of the piezo tuner.

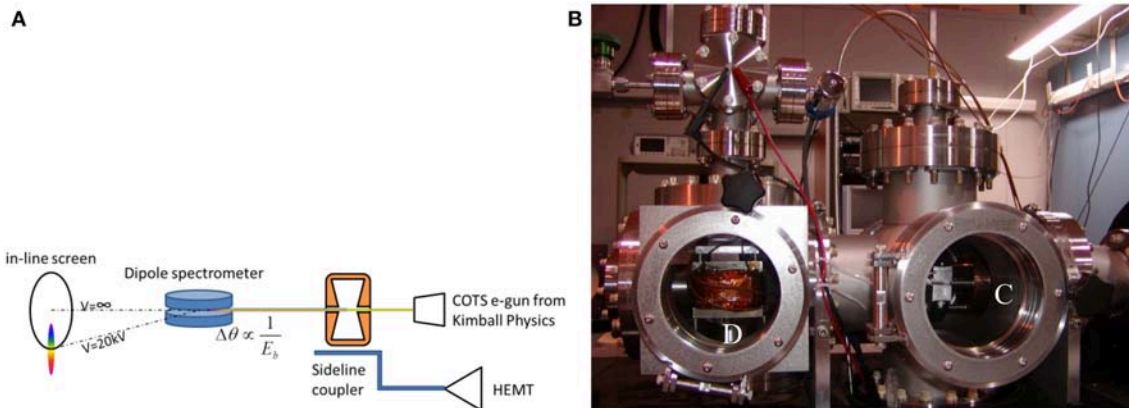


FIGURE 11 | (A) Schematic of the single-cavity energy gain measurement. The deflection angle is inversely proportional to the momentum of the electron beam; **(B)** Photo of the single-cavity energy gain measurement with the cavity (labeled C) mounted immediately after the Kimball-Physics electron gun (to the right of this picture), and the dipole spectrometer (labeled D).

$$\theta = \theta_0 e^{-\frac{t}{\tau}} = \theta_0 e^{-\frac{t}{(mC_V R_T)}} \quad (7)$$

Based on the temperature decay after RF power was turned off (time period O), we calculated a decay time constant of $\tau = 13,064 \text{ seconds} = 3.629 \text{ hours}$. From the cavity mass and heat capacity, the calculated thermal resistance is $R_T = 28.78^\circ \text{C/W}$. From Equation 6, we calculated the average RF power deposited into the cavity from the rate of temperature rise and the thermal resistance. During the two time periods labeled I and II in **Figure 9**, the HEMT duty factor was 5% (region I) and 10% (region II), and the calculated average power delivered to the cavity was 8.5 W (region I) and 17 W (region II). These average power measurements translate into a peak value of absorbed RF power of 170 W for both regions.

Temperature-Dependent Frequency Shift

As described above, the space-borne accelerator will not have temperature stabilization, so all cavities must be maintained at

the same frequency with active frequency control. We measured the resonant frequency of a C-band cavity as it was powered with 170 W at 5 and 10% duty factors without water cooling. As the cavity temperature rose by 65°C in 48 min, its resonant frequency decreased by 5.7 MHz ($-88 \text{ kHz}/^\circ \text{C}$) at an average rate of 2 kHz/s. To compensate for the temperature-induced frequency shifts, we designed a piezo tuner to be inserted into the cavity which would reduce the cavity inductance and thus shift the cavity resonant frequency to a higher frequency. The measured and CST-modeled cavity resonant frequency shifts vs. piezo tuner displacement in the cavity is shown in **Figure 10**. The range of piezo movement needed to compensate for the 5.7 MHz frequency shift due to the cavity temperature rise of 65°C is $<4 \text{ mm}$. We have found that both copper and aluminum make good material for the tuner as they preserve 99% of the cavity quality Q at the largest tuner displacement. In space, the accelerator will be mounted on a temperature-controlled surface and operated at approximately the same location in its orbit such

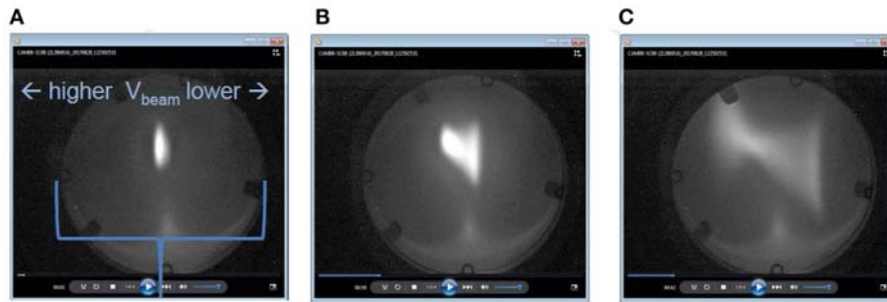


FIGURE 12 | (A) Position of the incoming electron beam on the screen with cavity RF off; **(B)** Energy-modulated electron beam with weak energy modulation at low RF power; and **(C)** Energy-modulated beam with strong energy modulation at medium RF power.

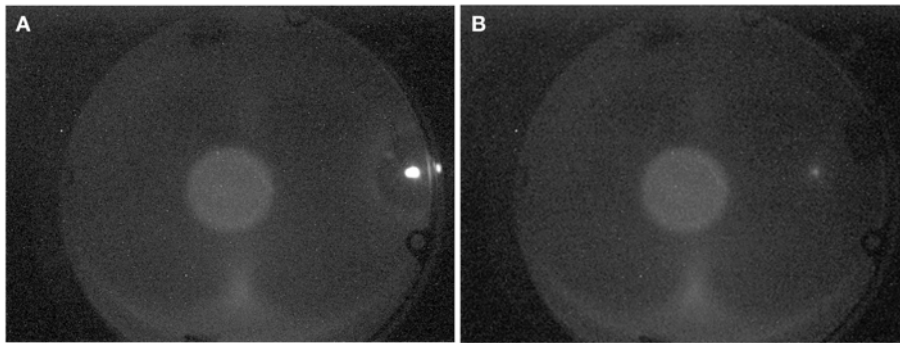


FIGURE 13 | (A) Position of the incoming electron beam with the cavity RF power turned off; **(B)** Position of the beam with cavity power turned on (the beam intensity was dimmer because only a small fraction of the incoming electrons were accelerated to the maximum energy).

that the temperature of the accelerator during operation will be between 15° and 25°C.

Energy Gain Measurement

The single-cavity energy gain was measured using the experimental setup illustrated in **Figure 11A**. A low-current 20-kV DC beam from a commercial Kimball Physics electron gun was passed through a prototype C-band cavity (labeled C in **Figure 11B**) powered with a single HEMT at various power levels. The deflection of the beam by a fixed-field dipole spectrometer (labeled D in **Figure 11B**) is proportional to the beam momentum, allowing measurement of the energy gain delivered to the beam.

With low-power RF delivered to the cavity, the DC electron beam experienced weak bidirectional energy shifts, i.e., energy modulations, and the energy gain was measured at the maximum energy of the energy-modulated beam on the screen. **Figures 12A–C** show images of the electron beams with no energy modulation (a), weak energy modulation with low cavity RF power (b) and strong energy modulations with medium RF power (c).

At even higher RF power (and thus higher energy gains), we were able to detect a well-defined beam on the screen corresponding to those electrons that are accelerated at the peak of the RF field. Since only a small fraction of the incoming electrons are at the peak of the RF field, the maximum-energy

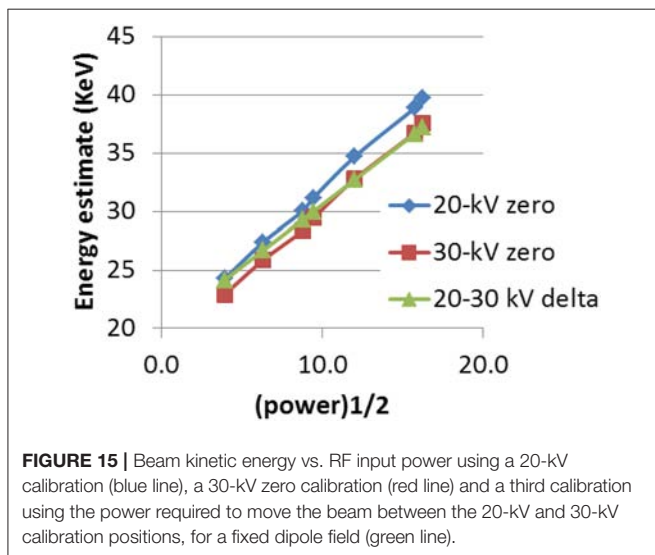
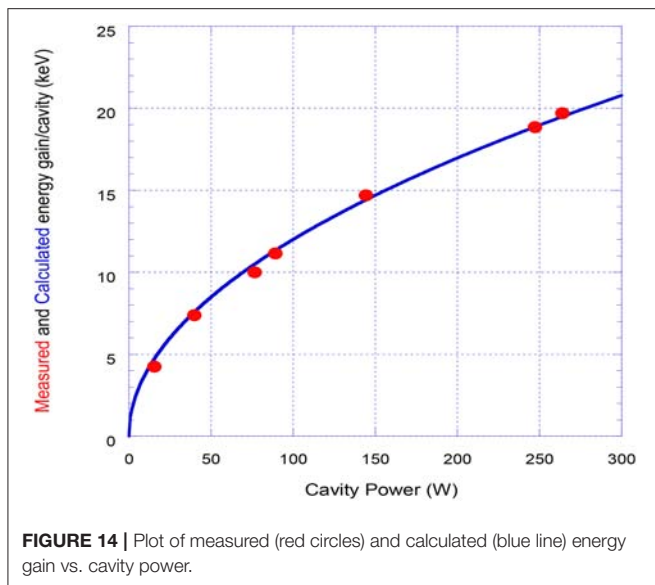
beam spot (**Figure 13B**) is dimmer than the incoming DC beam spot (**Figure 13A**).

With the dipole turned off (and degaussed) and the cavity RF off, we established the “zero” position of the incoming DC beam on the screen. This zero position was approximately centered on the round image of the faint cathode glow on the screen. Then, with RF power to the cavity still off, we adjusted the dipole field such that the incoming electron beam at 20 kV (called the 20-kV zero), or at 30 kV (called the 30-kV zero), impacted at the edge of the screen (**Figure 13A**). This calibrated the combination of dipole field and drift distance to displacement on the screen for a known beam energy. Next, as we increased the cavity RF power, the beam moved toward the center (**Figure 13B**) and the beam’s angular movement was used to determine the energy gain provided by the cavity.

The energy gain in an RF cavity is a product of the accelerating field and the transit-time corrected cavity length. Since the accelerating field is proportional to the square root of the RF power, the energy gain can be expressed by Equation 8.

$$\Delta W = E_0 L_{cav} = \sqrt{P_{cav} L_{cav} R_s} \quad (8)$$

Figure 14 shows the measured (red circles) and calculated (blue line) energy gain as a function of the cavity power. While the HEMT was operating at ~530 W, the cable connecting it to the cavity (inside a shielded enclosure) imposed ~3 dB attenuation,



limiting our maximum power to the cavity to 264 W. From the fit of the measured energy gain vs. the cavity power, we estimated the product of the shunt impedance per unit length (R_s) and the effective length (L_{cav}) of our C-band cavity to be $1.6 \pm 0.1 \text{ M}\Omega$.

We experimentally measured the beam energy gain as a function of RF power levels via three different methods: (1) By measuring the deflection of the incoming 20-kV beam from the screen edge as a function of cavity power (the 20-kV-zero method); (2) by measuring the deflection of the incoming 30-kV beam from the screen edge as a function of cavity power (the 30-kV-zero method); and (3) by measuring the cavity RF power needed to move the 20-kV beam from its zero position to the position of the 30-kV beam (the 20-30-kV-delta method). The beam kinetic energy gains measured by these

three methods are plotted vs. the square root of cavity power in **Figure 15**.

CONCLUSION AND FUTURE PLAN

We have demonstrated key aspects of technology required for the development of space-borne electron linacs. These include RF power source characterizations, tuner design, and beam acceleration with energy gain measurements. These initial test results show that HEMTs, operating with 50 V DC power supplies, can deliver sufficient RF power to individual accelerator cavities to provide energy gain of 20 keV per cavity. Raising the energy of the electron beam to 1 MeV will require approximately fifty of these C-band cavities, with each cavity powered by its own HEMTs and operated without water cooling. Compared to traditional klystron-based designs, the HEMT-powered linac design is more compact, efficient and suitable for space missions. It also avoids the use of high-voltage klystrons and associated power supplies which have been the single-point failures of terrestrial RF linacs. Currently, our team is concentrating effort on testing a multi-cavity prototype with the goal of accelerating the electron beam continuously in these cavities to a higher beam energy. The prototype will make use of an improved cavity design and an RF system that mimics a flight-appropriate system as closely as possible. As these cavities are physically independent, we may explore the possibility of using the first cavity as the buncher cavity, i.e., the first cavity will be used to modulate the energy of the incoming DC electron beam. The energy-modulated electrons will form short bunches of electrons at the cavity resonant frequency after drifting a short distance and these electron bunches will be captured and accelerated in the subsequent RF cavities. The phase and amplitude of the RF cavities can be independently adjusted to improve the fraction of electrons captured by the cavities. Finally, the multi-cavity prototype will allow exploration of various low-level RF control algorithms for maintaining cell-to-cell frequency and phase stabilization.

AUTHOR CONTRIBUTIONS

JL provided the physics design, performed all experimental work and analyzed the data. CB provided engineering designs, set up the beam experiment and assisted with the experimental work. BC was responsible for the first iteration of the CONNEX accelerator design. GD performed the initial HEMT tests and evaluation. MH and DP performed most of the HEMT testing and provided engineering support to the experiments. SS developed the engineering model of the CONNEX MeV space-borne accelerator. DN originated the concept, provided technical leadership, interpreted the data and prepared the manuscript with help from the co-authors.

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Effect of Field-Line Curvature on the Ionospheric Accessibility of Relativistic Electron Beam Experiments

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Magnetosphere-ionosphere coupling is a particularly important process that regulates and controls magnetospheric dynamics such as storms and substorms. However, in order to understand magnetosphere-ionosphere coupling it is necessary to understand how regions of the magnetosphere are connected to the ionosphere. It has been proposed that this connection may be established by firing electron beams from satellites that can reach an ionospheric footpoint creating detectable emissions. This type of experiment would greatly aid in identifying the relationship between convection processes in the magnetotail and the ionosphere and how the plasma sheet current layer evolves during the growth phase preceding substorms. For practical purposes, the use of relativistic electron beams with kinetic energy on the order of 1 MeV would be ideal for detectability. However, Porazik et al. (2014) has shown that, for relativistic particles, higher order terms of the magnetic moment are necessary for consideration of the ionospheric accessibility of the beams. These higher order terms are related to gradients and curvature in the magnetic field and are typically unimportant unless the beam is injected along the magnetic field direction, such that the zero order magnetic moment is small. In this article, we address two important consequences related to these higher order terms. First, we investigate the consequences for satellites positioned in regions subject to magnetotail stretching and demonstrate systematically how curvature affects accessibility. We find that curvature can reduce accessibility for beams injected from the current sheet, but can increase accessibility for beams injected just above the current sheet. Second, we investigate how detectability of ionospheric precipitation of variable energy field-aligned electron beams could be used as a constraint on field-line curvature, which would be valuable for field-line reconstruction and/or stability analysis.

Keywords: beam injection from space, field-line mapping, accessibility, loss cone, field-line curvature, energy-variable accelerator

1. INTRODUCTION

Plasma sheet transport is primarily driven by coupling of the magnetosphere and solar wind and has distinctly different behavior based on the orientation of the interplanetary magnetic field (IMF) with respect to the earth's dipole field (Wing et al., 2014, and references therein). Under northward IMF conditions, plasma convection weakens and transport may be dominated by turbulent flows (Borovsky and Funsten, 2003; Wang et al., 2010; Merkin et al., 2013). Under southward conditions, convection is stronger and may involve localized transient flows moving earthward (Angelopoulos et al., 1992; Sergeev, 2005; Birn et al., 2011; Wiltberger et al., 2015). Although transport of flux from the dayside to the nightside can be steady under southward IMF conditions, return of the flux to the dayside can be inhibited leading to the storage of flux in the plasma sheet and eventual release through substorms (Akasofu, 1964).

Observations in the ionosphere can provide insight into convection processes in the magnetotail (Sergeev, 2005; Bristow, 2008; Nishimura et al., 2010). The electric field responsible for magnetotail convection maps into the ionosphere (Ridley et al., 1998; Ruohoniemi and Baker, 1998) and auroral displays result in regions where flows twist magnetic fields, leading to field-aligned currents and electron precipitation. Diffuse particle precipitation detected by low altitude satellites can also provide a global picture of the plasma populations in the magnetotail (Wing and Newell, 1998; Wing et al., 2005; Wing and Johnson, 2009). However, connecting these ionospheric observations with magnetotail processes is complicated because the magnetic field mapping is not known precisely (Willis et al., 1997a,b). In order to understand the causal relationship between ionospheric observations and events/populations in the magnetosphere, it is necessary to map field-lines in the magnetosphere to their ionospheric footpoints. The use of empirical or MHD modeling techniques has made it possible to infer the mechanisms behind ionospheric observations; however, these results still involve significant uncertainty.

Having satellites configured with an electron beam generator is a promising method to map regions in the magnetosphere to the ionosphere by firing electrons into the loss cone and observing the precipitation from the ground (see Sanchez et al. in review). Experiments involving the artificial injection of electrons along magnetic field-lines in the magnetosphere has already shown feasibility of detecting electron beams (Winckler, 1980). However, it has been theorized (Neubert and Banks, 1992) that relativistic electron beams would be more stable than the beams used in these experiments (which had energies up to 40 keV) due to the higher relativistic mass and lower beam density. It is also suggested by simple linear analysis that relativistic beams traveling through the magnetosphere are stable to two-stream instabilities (Galvez and Borovsky, 1988), and relativistic beams entering the ionosphere are stable to resistive hose, ion hose, and filamentation instabilities (Gilchrist et al., 2001). Nevertheless, relativistic beams do come with their own issues, as discussed by Porazik et al. (2014), due to the fact that the first adiabatic invariant will not necessarily be conserved to zeroth order. Using a second order asymptotic expansion derived by Gardner (1966),

it was shown that the dependence on field-line curvature in the higher order terms of μ has a substantial effect on the loss cone. There are two important consequences of this fact that we discuss in this paper.

First, the fact that the loss cone is reduced by increasing field-line curvature is highly relevant in the case of satellites positioned near the equatorial plane at midnight local time. The magnetotail stretches during times of increased activity, which may cause the field-line curvature at the position of the satellite to increase significantly. However, the activity in these regions is relevant to understanding the magnetosphere-ionosphere connection. Therefore, it is useful to systematically examine how field-line curvature affects ionospheric accessibility in the regions under consideration in order to determine magnetotail configurations that would permit the technique to be used successfully without a significant reduction of beam precipitation due to curvature effects.

A second important consequence of this study is the possibility to infer field-line curvature by varying the energy of the beam. The curvature is an important variable that describes magnetotail stretching and current sheet thickness. As such, it would be particularly useful for considerations of stability of the magnetotail to ballooning instability and/or reconnection. For example, if the magnetic field curvature is known, it would provide a significant constraint of magnetotail equilibria and therefore could potentially be used to constrain equilibrium models used for stability analysis (Cheng, 1995; Cheng and Zaharia, 2004).

A relatively simple threshold condition relating beam energy and curvature can be obtained when the accelerator is aimed in the direction of the magnetic field. For field-aligned electrons, Gardner's formula takes the simple form:

$$\mu = \bar{\mu} \rho^2 \kappa^2 \quad (1)$$

where $\rho = \gamma mv/qB$, $\bar{\mu} = \gamma mv^2/2B$, and κ is the magnitude of the field-line curvature, given by $\vec{\kappa} = (\hat{b} \cdot \nabla) \hat{b}$ where $\hat{b} = \vec{B}/B$. Field-aligned electron beams are therefore expected to precipitate if

$$\rho^2 \kappa^2 < B/B_i \quad (2)$$

where B_i is the magnitude of the field at the ionospheric footpoint. These beams can therefore be used to obtain information about the field-line curvature at the launch position. If precipitation of the beam is observed, then it must be true that

$$R_c > \rho \sqrt{\frac{B_i}{B}} \quad (3)$$

where $R_c = 1/\kappa$. Although obtaining a lower bound on the radius of curvature may be useful, this relationship seems to reveal an opportunity for directly measuring the curvature if the accelerator is capable of varying the energy of the beam. If all of the particles are fired exactly in the direction of the field, then one would only need to increase the energy until precipitation is no longer observed, indicating that the above inequality is no

longer satisfied and $R_c = \rho \sqrt{B_i/B}$. However, no accelerator will be capable of firing every electron exactly in the direction of the field. It must be shown that, for an electron beam aimed in the direction of $\vec{B}$ and having some pitch angle spread Δ , the fraction of particles simultaneously fired that are in the loss cone will significantly decrease at the critical energy where $\rho = R_c \sqrt{B/B_i}$ in order to fully validate this concept.

2. METHODOLOGY

At a given energy and initial launch position, we will express the initial velocity of electrons by the angles (ϕ, α) , which are defined by:

$$\begin{aligned}\tan(\phi) &= \frac{\vec{v} \cdot \hat{N}}{\vec{v} \cdot \hat{N}_b} \\ \cos(\alpha) &= \frac{\vec{v} \cdot \vec{B}}{vB}\end{aligned}\quad (4)$$

where $\hat{N}$ and $\hat{N}_b$ are the normal and bi-normal vectors of the field line at the launch point, respectively. Note that ϕ and α are merely the azimuthal and lateral angles in conventional spherical coordinates where the z-axis is aligned with $\vec{B}$. Given an initial launch position $\vec{x}$, a beam energy E , and initial velocity defined by the pair of angles (ϕ, α) , we denote the value of the magnetic moment for electrons given by Gardner's formula by $\mu(\vec{x}, E, \phi, \alpha)$. We then define the loss cone as the set:

$$LC(\vec{x}, E) = \left\{ (\phi, \alpha) \mid \mu(\vec{x}, E, \phi, \alpha) \leq \gamma m v^2 / 2B_i \right\} \quad (5)$$

If μ is conserved to second order, ionospheric accessibility at a given energy and launch position can be expressed through the surface area of the loss cone in velocity space, given by

$$A(\vec{x}, E) = \int_{LC(\vec{x}, E)} \sin(\alpha) d\alpha d\phi \quad (6)$$

see the **Appendix** for a detailed explanation of how A may be computed most efficiently. If μ is not conserved to second order, then A can not be expected to represent ionospheric accessibility. This fact is especially important in the case where A is not monotonic along field-lines. In many of these cases, A will change sharply along the field-line, which conflicts with the assumption that μ is conserved to second order. For this reason, we will consider an alternative accessibility metric to A :

$$A^*(\vec{x}, E) = \min \left\{ A(\vec{y}, E) \mid \vec{y} \in X \right\} \quad (7)$$

where X is the set of points containing the point $\vec{x}$ that all lie on the same field-line and lie between the point $\vec{x}$ and the ionospheric footprint.

In order to validate the idea of using variable energy accelerators to measure field-line curvature, we must investigate the accessibility of field-aligned beams near the critical energy, defined by:

$$E_c = mc^2 \left(\sqrt{1 + \frac{\Omega^2 R_c^2}{c^2} \frac{B}{B_i}} - 1 \right) \quad (8)$$

where $\Omega = qB/m$. Note that Ω is the non-relativistic gyrofrequency and the actual gyrofrequency is Ω/γ . E_c is the energy where $\rho = R_c \sqrt{B/B_i}$, and is therefore the threshold energy above which electrons having initial velocity exactly aligned with $\vec{B}$ will not precipitate in the ionosphere. For the sake of simplicity, we will suppose that the instantaneous density of the beam is normally distributed in α with standard deviation σ :

$$n(\alpha) = \frac{\exp \frac{-\alpha^2}{2\sigma^2}}{\sqrt{2\pi^3} \sigma \operatorname{erf}(\frac{\Delta}{\sigma\sqrt{2}})} \quad (9)$$

where Δ is the pitch angle spread of the beam. One can check that this distribution satisfies the normalization condition:

$$\int_0^{2\pi} \int_0^\Delta n(\alpha) d\alpha d\phi = 1 \quad (10)$$

Given a beam energy and initial launch position, the fraction of particles instantaneously fired from the beam that are in the loss cone can be determined simply by drawing a sufficiently large number of pairs (ϕ, α) from this distribution and computing the fraction of those pairs that are in the set LC .

3. RESULTS

In order to systematically investigate the consequences of tail stretching for beam accessibility, we calculate A on the midnight meridional plane using the Tsytganenko 1989 model (Tsytganenko, 1989) with K_p ranging between 1 and 7, where we assume a 1 MeV beam (see **Figure 1**). Field-lines are also shown so that the changes in the field-line curvature can be seen visually. As K_p increases, it is visibly apparent that the field-line curvature increases on the equatorial plane and decreases away from the equatorial plane. The dark regions indicate regions of relatively small A , which implies low beam accessibility. For $K_p = 1$ and $K_p = 2$, A is smallest in regions where field-line curvature is greatest. However, for $K_p > 2$, this trend is broken for $X < -7R_E$. In these cases, A is seen to be large on the equatorial plane, and essentially vanishes in regions immediately above and below the equatorial plane. This reduction in accessibility is due to terms in Gardner's formula that depend on $\partial\kappa/\partial n$ or $\partial\kappa/\partial s$, where $\partial/\partial n = \hat{N} \cdot \nabla$ and $\partial/\partial s = \hat{b} \cdot \nabla$. If the higher order derivatives of the curves are large, then these terms increase μ so that A is significantly reduced. However, for $|Z| > 1R_E$, A consistently increases as K_p increases due to the reduction of field-line curvature in these regions.

It should be noted that although the equatorial μ in these cases can be consistent with ionospheric precipitation, it is unlikely in this case that μ is actually conserved through the dark regions just above the equatorial regions where A vanishes. As mentioned previously, A is only a reflection of beam accessibility if μ is conserved to second order. However, high gradients in A conflict with the assumption that μ is conserved, implying that electron beams are not generally accessible to the ionosphere when launched from any region equatorward of the region of low accessibility. In other words, if a particle were to move along the field line from the equatorial plane into a dark region, there

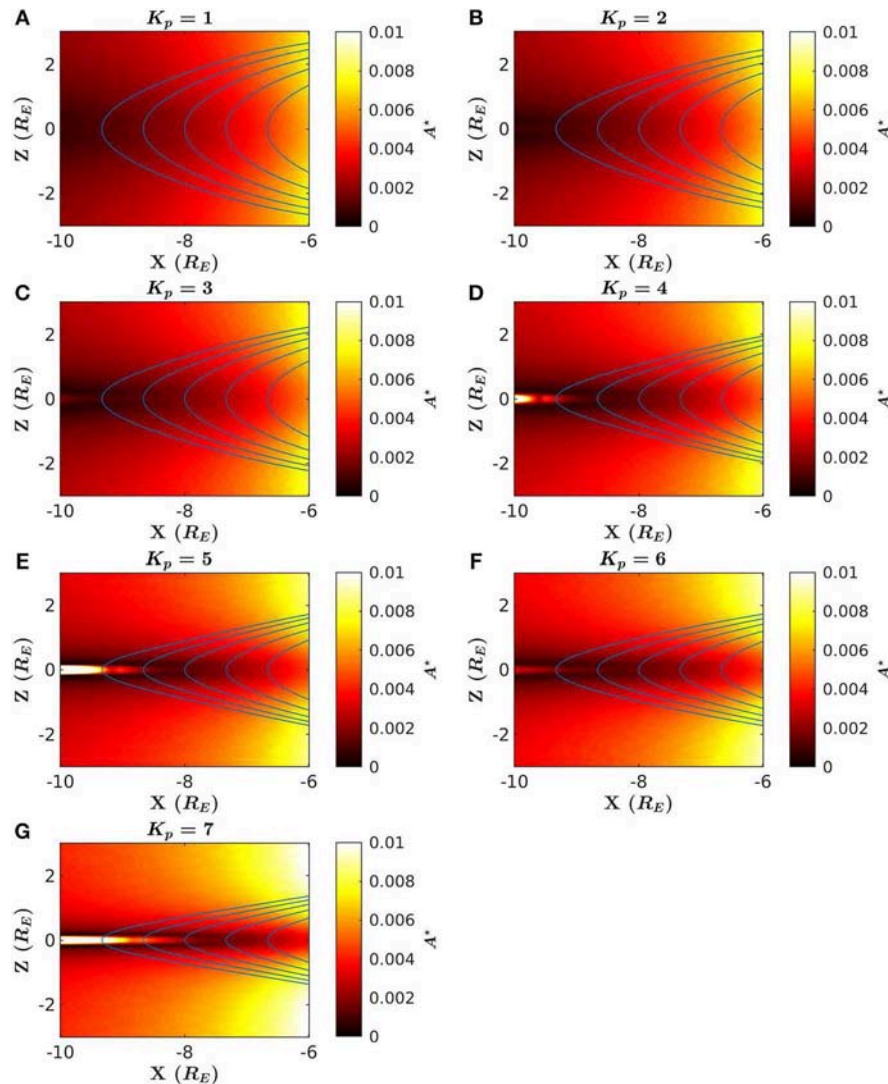


FIGURE 1 | A plotted over midnight local time for various K_p assuming a beam energy of 1 MeV: (A) $K_p = 1$, (B) $K_p = 2$, (C) $K_p = 3$, (D) $K_p = 4$, (E) $K_p = 5$, (F) $K_p = 6$, (G) $K_p = 7$. Blue curves represent field-lines that intersect chosen values of X on the equatorial plane. Darker and lighter regions indicate regions with smaller and greater values of A , respectively. Note that A is not expected to reflect accessibility at the equatorial plane for high K_p .

would be no path of accessibility from that dark region to the ionosphere that conserves magnetic moment. For this reason, A^* is a better reflection of beam accessibility (see **Figure 2**). From this figure, it is more clearly seen that accessibility at the equatorial plane decreases as K_p increases. It should be noted that there may be paths of accessibility for particles for which the magnetic moment is not conserved, but such precipitation cannot be predicted reliably due to the chaotic nature of orbits in the plasma sheet (Chen et al., 1990).

If μ is conserved to second order, then a beam fired with pitch angle α that precipitates in the northern hemisphere is also expected to precipitate in the southern hemisphere if the pitch angle were instead chosen to be $\pi - \alpha$. However, in the same way that μ is not expected to be conserved for electrons fired from the equatorial plane, μ is likewise not expected to be conserved for electrons passing through the equatorial plane. Therefore, although electron beams away from the equatorial

plane are reliably accessible to at least one of the two hemispheres, the beams cannot be simultaneously accessible to both the northern and southern hemispheres when the tail is stretched. This consequence must be taken into consideration if ground observers are not stationed at locations both in the northern and southern hemispheres. It should also be noted that this limitation is not seen for non-relativistic beam experiments (Winckler, 1980) or for experiments involving satellites in geosynchronous orbit.

Figure 3 shows the fraction of particles fired into the loss cone in the case of a field-aligned beam with realistic beam spread $\Delta = 0.005$ (see Sanchez et al. in review), with $\sigma = \Delta/3$, and for energies above and below the E_c . All other parameters are chosen to correspond with a satellite positioned on local time midnight, at the equatorial plane, and at a radial distance of $8 R_E$ with $K_p = 1$. It is seen that the fraction of particles in the loss cone decreases significantly at E_c . If μ is conserved to second order,

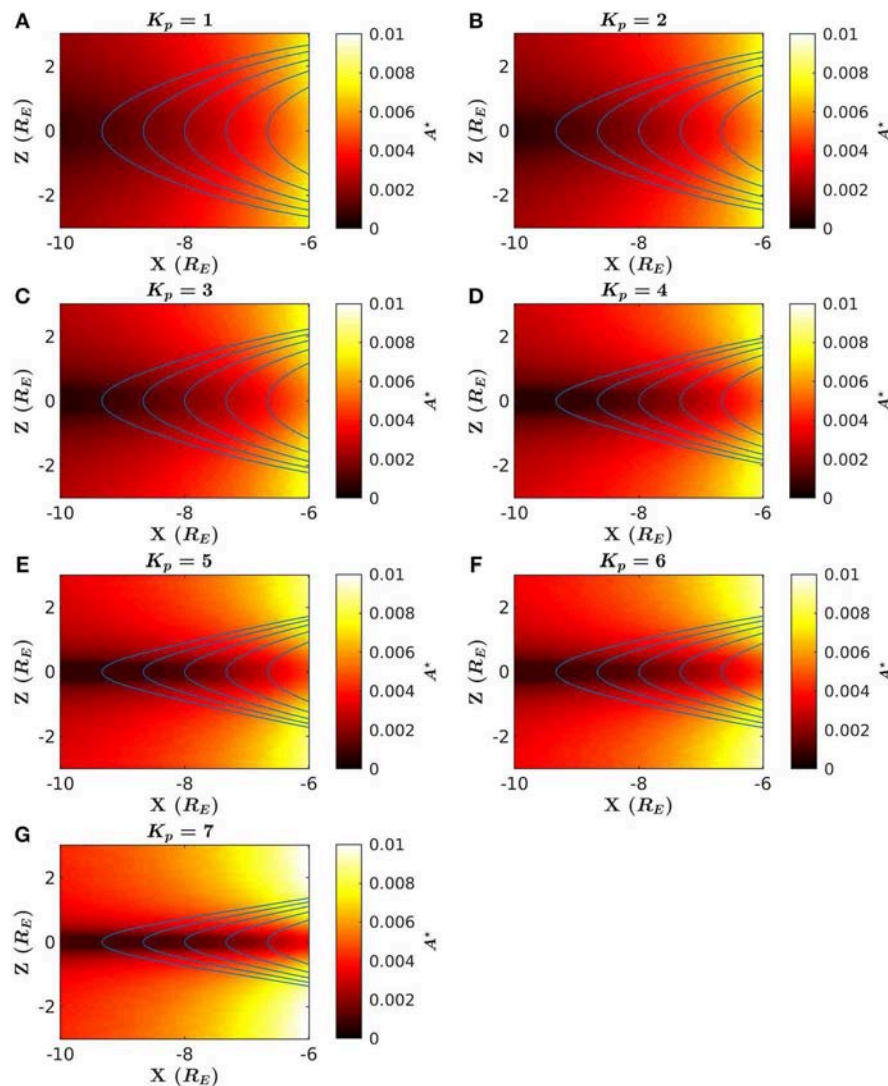


FIGURE 2 | A^* plotted over midnight local time for various K_p assuming a beam energy of 1 MeV: **(A)** $K_p = 1$, **(B)** $K_p = 2$, **(C)** $K_p = 3$, **(D)** $K_p = 4$, **(E)** $K_p = 5$, **(F)** $K_p = 6$, **(G)** $K_p = 7$. Blue curves represent field-lines that intersect chosen values of X on the equatorial plane. Darker and lighter regions indicate regions with smaller and greater values of A^* , respectively.

then it follows that the fraction of particles that precipitate in the ionosphere should be expected to decrease significantly at this threshold.

4. CONCLUSIONS

The results above suggest that as the magnetotail is stretched and magnetic curvature increases there will be a reduction of accessibility in the equatorial plane $|Z| < 1R_E$ and an increase of accessibility in regions where $|Z| > 1R_E$ where μ is expected to be conserved. In these regions, increasing K_p corresponds with increasing A^* in all cases, indicating that ionospheric accessibility is increased when the tail is stretched and field-line curvature is reduced. These results show that tail stretching may create difficulty in performing electron beam experiments if the satellite is positioned on the equatorial plane. However, because this

issue is not seen for launch positions where $|Z| > 1R_E$, it may still be possible to use relativistic beams to map field-lines from the ionosphere to locations near the plasma sheet. This result validates one aspect of the idea that electron beams may be used to study the relationship between plasma sheet dynamics and observations in the ionosphere. Alternatively, our results generally suggest that relativistic electron beams could be used to map field lines within geosynchronous orbit for a wide range of geomagnetic conditions.

For field-aligned beams, we have shown that the fraction of particles that precipitate in the ionosphere decreases sharply around E_c . If the accelerator were capable of varying the beam energy from below E_c , ground observers would be able to detect the change in precipitation at E_c and would therefore be able to infer the field-line curvature (see Powis et al., in review, for details surrounding ground observations of electron beam experiments

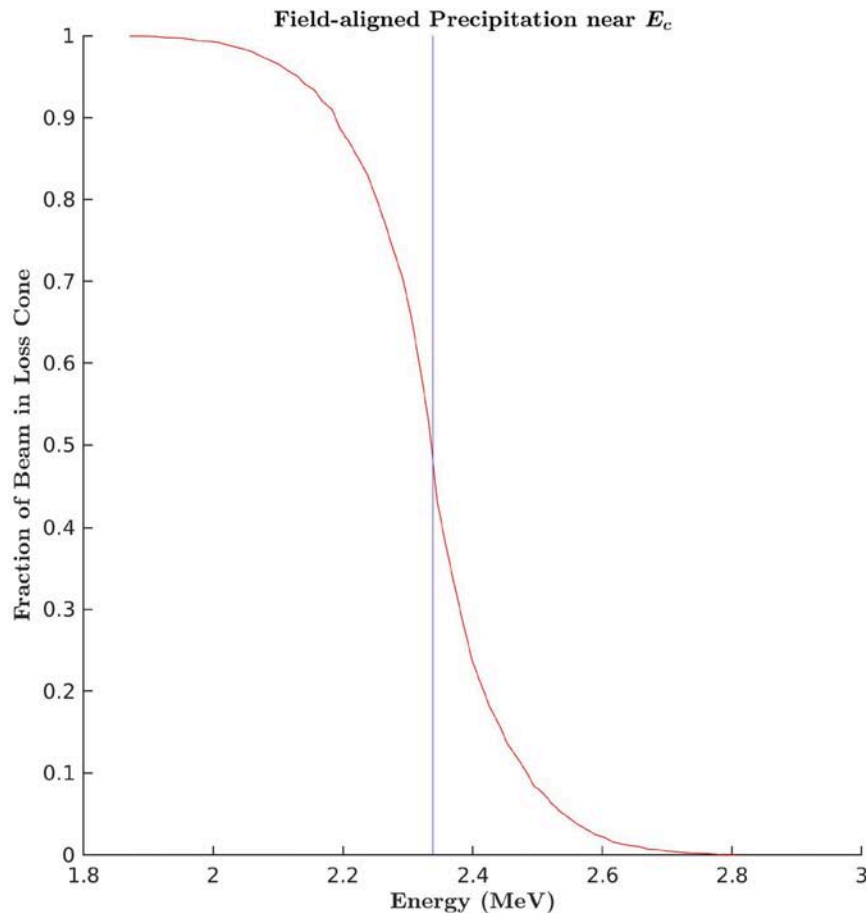


FIGURE 3 | The fraction of particles fired into the loss cone plotted against beam energy in the case of a field-aligned beam with $\Delta = 0.005$, $\sigma = \Delta/3$. E_c is represented by the blue line.

in general). The ability to directly measure aspects of the field-line geometry would be a substantial aid in constraining models of the magnetic field of Earth's inner magnetosphere. In addition to measuring curvature, the arc length of field-lines between the launch and precipitation points may be directly measured using low energy keV field-aligned beams. In this case, μ is well approximated by the zeroth order term, so $v_{\parallel} \approx v$. The total distance that the electrons travel is therefore equal to the arc-length of the field-line. The ability to measure both the arc length and curvature of field-lines in the magnetosphere using a single satellite is a very attractive opportunity, and our results demonstrate that the issue of using this kind of information to properly constrain models is a topic worthy of extensive study (see Willard et al. in review).

These results do not provide insight into situations involving magnetic fields that are not well described by the Tsyganenko 1989 model, which could occur in a turbulent or reconnecting plasma sheet where small-scale curvature could be found. However, we expect that such configurations offer less accessibility than that described by the Tsyganenko model. These results also do not provide insight into how the accessibility may be affected by waves found in the magnetosphere or created by

the beam itself, which may have the effect of scattering electrons out of the loss cone even if they are initially fired into it. However, calculations carried out by Glauert and Horne (2005) seem to show that pitch angle scattering based on typically observed wave amplitudes is not significant for relativistic electrons with small initial pitch angle that interact with whistler mode, EMIC, and Z mode waves. In these cases, no bounce-averaged diffusion coefficient was determined for electrons with small initial pitch angle since these electrons were not immediately scattered out of the loss cone, which indicates that wave particle interaction would not be a significant effect for electron beams fired with small pitch angle. Nevertheless, Glauert and Horne's calculations do not show that this is generally the case for beams fired into the loss cone, and further investigation of beam stability and wave saturation is required in order to completely rule out the significance of this effect on the overall accessibility.

AUTHOR CONTRIBUTIONS

JW wrote the manuscript and made the calculations with help from JJ, in consultation with ES, JS, AP, and IK. JJ devised the

project. All authors discussed the results prior to and during the preparation of the manuscript.

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A. APPENDIX

The angular coordinates used to formalize the concept of loss cone surface area, although useful for the purpose of describing the concept, are not the preferred coordinates to use when calculating A . For this purpose, it is convenient to apply a coordinate transformation $(\lambda, \beta) = T(\phi, \alpha)$, where the initial velocity corresponding to the angles (λ, β) satisfy the relations:

$$\begin{aligned}\tan(\lambda) &= \frac{\vec{v} \cdot \hat{e}_1}{\vec{v} \cdot \hat{e}_2} \\ \cos(\beta) &= \frac{\vec{v} \cdot \vec{v}_0}{v^2}\end{aligned}\quad (\text{A1})$$

where $\vec{v}_0 = v\hat{e}_1 \times \hat{e}_2$ is the initial velocity of the trajectory having $\mu = 0$. Note that, like the previous coordinate system, this is nothing more than a conventional spherical coordinate system where the z -axis is aligned with a particular direction. In this case, λ and β are azimuthal and lateral angles for a spherical coordinate system with the z -axis aligned with $\vec{v}_0$. As described by Porazik et al. (2014), $\vec{v}_0$ is not exactly aligned with the magnetic field and will always have a perpendicular component that is aligned with the drift. The transformation T is therefore simply a rotation of the system about the normal axis $\hat{N}$ by an angle α_0 such that $\cos(\alpha_0) = \vec{v}_0 \cdot \vec{B}/vB$. Setting $\hat{v}_0 \cdot \hat{N} = 0$ in Gardner's formula, α_0 can be easily computed by solving

$$\begin{aligned}(K_1^2 - K_3 + K_4)w^4 + (K_2 - 2K_1)w^3 + (1 + K_3 - 2K_1^2)w^2 \\ + 2K_1w + K_1^2 = 0\end{aligned}\quad (\text{A2})$$

where $w = \sin(\alpha_0)$ and

$$\begin{aligned}K_1 &= \rho\kappa \\ K_2 &= \frac{\rho}{B} \frac{\partial B}{\partial n} \\ K_3 &= \rho^2 \left(\frac{1}{2r^2} \frac{B_r^2}{B^2} + \frac{B_r}{2rB^2} \frac{\partial B}{\partial s} - \frac{\kappa B_z}{4rB} + \frac{1}{8} \left(\frac{1}{B} \frac{\partial B}{\partial s} \right)^2 - \frac{11}{4} \kappa^2 \right. \\ &\quad \left. + \frac{9\kappa}{2B} \frac{\partial B}{\partial n} - \frac{7}{4} \frac{\partial \kappa}{\partial n} \right) \\ K_4 &= \rho^2 \left(\frac{B_r^2}{8r^2B^2} - \frac{B_r}{8rB^2} \frac{\partial B}{\partial s} - \frac{B_z}{8rB^2} \frac{\partial B}{\partial n} - \frac{5}{32} \left(\frac{1}{B} \frac{\partial B}{\partial s} \right)^2 \right. \\ &\quad \left. + \frac{15}{8} \left(\frac{1}{B} \frac{\partial B}{\partial n} \right)^2 - \frac{5}{8} \frac{1}{B} \frac{\partial^2 B}{\partial n^2} \right)\end{aligned}\quad (\text{A3})$$

The precise choice of unit vectors $\hat{e}_1$ and $\hat{e}_2$ is not important so long as $\hat{e}_1 \cdot \hat{e}_2 = 0$. This degree of ambiguity merely corresponds to a phase offset in the angle λ . For this analysis, we make the choice that $\lambda = 0$ should correspond to velocities where $\vec{v} \cdot \hat{N} = 0$ and $\vec{v} \cdot \vec{B} \times \nabla B > 0$.

In the original coordinate system, it was found by Porazik et al. that a given angle ϕ may correspond with two different points on the loss cone boundary. Conveniently, this is not a feature of these new coordinates, and one can check that every point on the loss cone boundary has a unique value of λ . This fact makes the integral A far more straight forward to calculate. In these new coordinates, the integral takes the form

$$A = \int_0^{2\pi} [1 - \cos(\beta_b(\lambda))] d\lambda \quad (\text{A4})$$

where $\mu = \gamma mv^2/2B_i$ when $(\lambda, \beta) = (\lambda, \beta_b(\lambda))$. The above expression can be very easily approximated with high precision using the trapezoidal rule.



Method for Approximating Field-Line Curves Using Ionospheric Observations of Energy-Variable Electron Beams Launched From Satellites

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Using electron beam accelerators attached to satellites in Earth orbit, it may be possible to measure arc length and curvature of field-lines in the inner magnetosphere if the accelerator is designed with the capability to vary the beam energy. In combination with additional information, these measurements would be very useful in modeling the magnetic field of the inner magnetosphere. For this purpose, a three step data assimilation modeling approach is discussed. The first step in the procedure would be to use prior information to obtain an initial forecast of the inner magnetosphere. Then, a family of curves would be defined that satisfies the observed geometric attributes measured by the experiments, and the prior forecast would then be used to optimize the curve with respect to the allowed degrees of freedom. Finally, this approximation of the field-line would be used to improve the initial forecast of the inner magnetosphere, resulting in a description of the system that is optimally consistent with both the prior information and the measured curvature and arc length. This article details the method by which a family of possible approximations of the field-line may be defined via a numerical procedure, which is central to the three step approach. This method serves effectively as a pre-conditioner for parameter estimation problems using field-line curvature and arc length measurements in combination with other measurements.

Keywords: field-line geometry, data assimilation, field-line approximation, beam injection from space, energy-variable accelerator

1. BACKGROUND

1.1. Motivation

Current accelerator technologies allow for the possibility of equipping small to medium satellites with lightweight electron beam accelerators. The scientific potential of such a setup is that an electron beam can be fired into the loss cone from somewhere in the inner magnetosphere and will end up in the ionosphere. Simulations have shown that electron beams fired into the ionosphere would result in observable precipitation (Marshall et al., 2014, 2019), and this could allow for the mapping of field-lines in the inner magnetosphere to their ionospheric foot-points at altitudes

ranging between 40 and 75 km above the Earth's surface (Marshall and Bortnik, 2018). Past studies suggest that utilizing both relativistic and non-relativistic electron beams in this way is possible and would provide a means of mapping field-lines at important locations in the inner magnetosphere to the ionosphere. Experiments involving the artificial injection of non-relativistic electron beams (having energies up to 40 keV) have shown that detecting these beams is feasible (Winckler, 1980), and relativistic electron beams are expected to be more stable due to higher relativistic mass and lower beam density (Neubert and Banks, 1992). Additionally, simple linear analysis suggests that relativistic beams traveling through the magnetosphere are stable to two-stream instabilities (Galvez and Borovsky, 1988), and are stable upon entering the ionosphere to resistive hose, ion hose, and filamentation instabilities (Gilchrist et al., 2001).

A necessary consideration for this endeavor is the magnetic moment μ of the electrons in the beam, as the beam will not precipitate unless $\mu < \gamma mv^2/2B_i$, where B_i is the magnitude of the field at the ionosphere. At midnight local time in the inner magnetosphere, we may be able to employ a second order asymptotic expansion of μ derived by Gardner (1966), which is valid under conditions discussed in the following section. If we assume that the beam is fired strictly in the direction of the magnetic field, the formula takes a very simple form:

$$\mu = \bar{\mu} \rho^2 \kappa^2 \quad (1)$$

where $\rho = \gamma mv/qB$, $\bar{\mu} = \gamma mv^2/2B$, and κ is the magnitude of the field-line curvature, given by $\vec{\kappa} = (\hat{b} \cdot \nabla)\hat{b}$ where $\hat{b} = \vec{B}/B$. Field-aligned electron beams are therefore expected to precipitate if

$$\rho^2 \kappa^2 < B/B_i \quad (2)$$

This relationship reveals an opportunity for obtaining significant information about the field-line geometry if the satellite is capable of varying the energy of the beam. In the case of low energy beams (on the order of 1 keV), μ is well-approximated by the zeroth order term, so $v_{\parallel} \approx v$. Since the total distance traveled by the particles must then be approximately equal to the arc length of the field-line between the launch and precipitation points, the arc length can be inferred by measuring the electron time of flight. If the energy of the beam is then increased, there may eventually come a critical point where observed precipitation is significantly reduced, indicating that the above inequality is no longer satisfied. Taking ρ_c to represent the corresponding value of ρ at this critical energy, the radius of curvature at the launch point is determined to be:

$$R_c = \rho_c \sqrt{\frac{B_i}{B}} \quad (3)$$

This concept has been expanded in greater detail by Willard et al. (2019).

The ability to measure both field-line curvature and arc length using a single satellite would significantly improve our ability to model the magnetic field of the inner magnetosphere.

However, the issue of how to properly constrain a model using field-line geometry has not been investigated thoroughly. In particular, using measurements of field-line arc length in the context of parameter estimation would seem to require a highly inefficient procedure. For a given choice of model parameters for the magnetic field, the error in the field-line arc length must be computed by tracing field-lines numerically and computing their arc lengths. This means that the standard approach to parameter estimation, where the error is minimized iteratively, would require solving a non-linear initial value problem once per iteration, substantially increasing computational complexity in comparison to typical parameter estimation problems.

The method presented in this paper is motivated by an alternative approach, which will require that there is a way to use the measurements of field-line geometry to approximate the field-line curves themselves. If reasonable approximations of field-lines are possible, then these approximations can be used to enable parameter estimation without requiring the complex calculation previously described. Instead, model parameters may be estimated to maximize the alignment of the field with the approximated curve, which is a far more straightforward task. An especially elegant case where this could be used is in the context of equilibrium models, where field-line curves are already used to establish boundary conditions (Cheng, 1995; Zaharia et al., 2004, 2005). No matter the approach used, it is certainly true that making predictions about the magnetic field configuration in the inner magnetosphere must require more than just arc length and curvature measurements, and this must also be true of any effort to approximate field-lines themselves. It is therefore useful to consider the method described in this paper in the context of a three step data assimilation procedure (see **Figure 1**), where field-line geometry measurements would be combined with prior information about the inner magnetosphere. Data assimilation methods have been used in geospace science for some time (Richmond, 1992; Schunk et al., 2004; Kondrashov et al., 2007; Merkin et al., 2016), with no shortage of models that may be used to describe geospace systems (Cheng, 1995; Lyon et al., 2004; Tóth et al., 2007; Janhunen et al., 2012). The first step of this procedure would be to use the prior information to obtain a naive forecast of the magnetic field using traditional parameter estimation techniques. Next, the field-line geometry would be used to infer the field-line curve itself, satisfying the measured geometric restrictions, and would be chosen to be as consistent as possible with the naive forecast. Finally, the field-line curve and the prior information would be used to obtain a complete forecast of the magnetic field, again using traditional parameter estimation techniques.

In the context of this kind of three step approach, it is clear that the approximation of field-lines is merely a means of transforming the measured geometric information into a form that can be more easily used to constrain magnetic field models using parameter estimation. Our method should therefore be thought of as a pre-conditioner for the original problem, provided that there is enough additional data available such that the problem can be solved in the first place. The purpose of this paper is to present a method by which a family of curves,

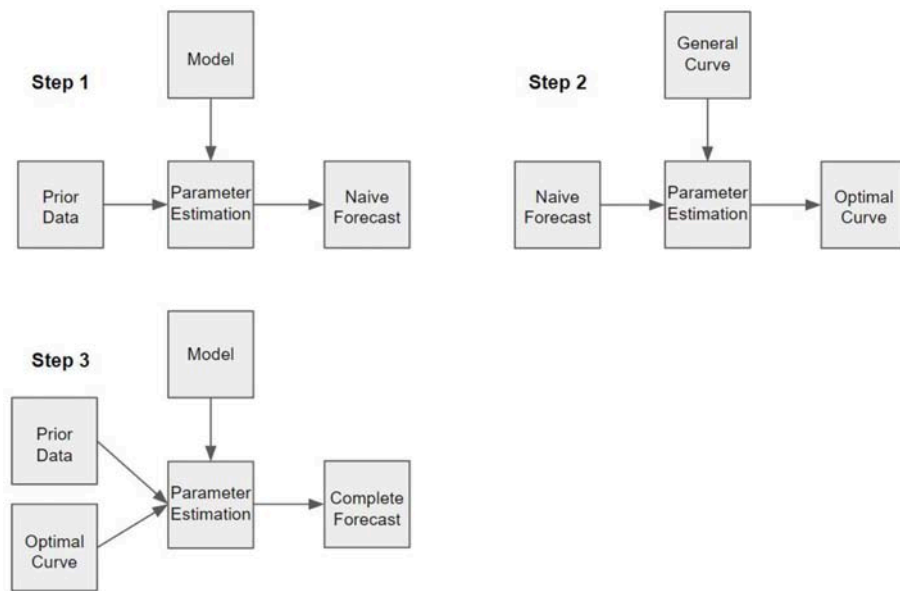


FIGURE 1 | Schematic of the three step data assimilation procedure. In each step, some model is being optimized with respect to provided information. In step one, the physical model is being optimized with respect to prior data to produce a naive forecast. In step two, a family of possible field-line curves is optimized with respect to the naive forecast to produce an optimal field-line curve, which is the best approximation of the field-line. In step three, both the prior data and the optimal curve are used to optimize the physical model and produce a complete forecast.

satisfying the measured geometric attributes, may be defined such that all remaining degrees of freedom are expressed in terms of a finite set of free parameters, since it is these free parameters that would need to be optimized in step two of the procedure (Figure 1) resulting in a unique approximation of the field-line.

1.2. Applicability

Since the scope of this method will be restricted only to situations involving data taken from energy-variable electron beam experiments, it is important to clarify the key assumptions upon which the inference of field-line arc length and curvature are based. Central to the former inference is the assumption that the path of the electron beam very closely approximates the field-line curve, so that the arc length is approximately equal to the total distance traveled by the electrons. For field-aligned electron beams, this can only be assumed if μ is dominated by the zeroth order term, which requires $\rho\kappa \ll 1$. We also must assume that the field-line geometry does not shift appreciably over the particle time of flight, which is on the order of seconds. Central to the later inference is the assumption that $B_\phi \approx 0$ and that the field can be well-approximated as axisymmetric to second order, since this is required in order for Gardner's formula to represent an adiabatic invariant. In part, the validity of these assumptions are energy dependent. For field-aligned particles, we can express the kinetic energy as a function of μ :

$$E(\mu) = mc^2 \left(\sqrt{1 + \frac{\Omega^2 R_c^2}{c^2} \frac{\mu}{\bar{\mu}}} - 1 \right) \quad (4)$$

where $\Omega = qB/m$. From this, we define three relevant energies:

$$\hat{E}_1 = mc^2 \left(\sqrt{1 + \frac{\Omega^2 R_c^2}{c^2} (0.1\%) - 1} \right) \quad (5a)$$

$$\hat{E}_2 = mc^2 \left(\sqrt{1 + \frac{B}{B_i} \frac{R_c^2 \Omega^2}{c^2} - 1} \right) \quad (5b)$$

$$\hat{E}_3 = mc^2 \left(\sqrt{1 + \frac{\Omega^2 L_\phi^2}{c^2} - 1} \right) \quad (5c)$$

$\hat{E}_1$ is the energy where $\mu/\bar{\mu} = 0.1\%$, and is therefore a reasonable maximum energy where the approximation $v_{\parallel} \approx v$ is valid. $\hat{E}_2$ is the critical energy where the inequality (2) is violated, and $\hat{E}_3$ is the energy where $\rho = L_\phi$, where L_ϕ is defined by

$$L_\phi^2 = r^2 \sin^2(\theta) \left(\frac{1}{B} \frac{\partial^2 B}{\partial \phi^2} \right)^{-1} \quad (6)$$

L_ϕ is a distance scale corresponding to the second order variation of B in ϕ . We must assume axisymmetry in order to employ Gardner's formula, so $\hat{E}_3$ is therefore the energy where Equation (1) should no longer be expected to apply to field-aligned particles. Note, however, that ρ is not necessarily equal to the gyroradius unless the particle motion becomes completely perpendicular to the direction of the field. It is understood then that $\hat{E}_3$ represents a conservative restriction on where validity in our assumptions is expected. Since the altitude of peak precipitation depends slightly on the energy of the beam, it is important to recognize that B_i is similarly dependent on energy.

We have neglected this dependence in the above definitions in light of recent simulations which have shown that electron beams having energies between 0.1 and 10 MeV will observably precipitate within a range of altitudes spanning roughly 35 km (Marshall and Bortnik, 2018), which does not correspond to a significant variation in B_l . In order to perform this experiment successfully, the satellite must be capable of varying the beam energy to reach both $\hat{E}_1$ and $\hat{E}_2$, the energy of the beam must be less than $\hat{E}_3$ everywhere on its path, and $\hat{E}_2$ must be less than $\hat{E}_3$ at the launch point.

Using T89, the Tsyganenko 1989 magnetic field model (Tsyganenko, 1989), we are able to show precisely how these energy criteria affect the scope of the method. In **Figure 2**, the thresholds $\hat{E}_1$, $\hat{E}_2$, and $\hat{E}_3$ are each calculated over midnight local time in the case where $K_p = 1$, where K_p is the global geomagnetic activity index. The color scale displays low energies using brighter colors, and high energies using darker colors. Along with these calculations, **Figure 2** also indicates the region on midnight local time where the criteria $10 \text{ keV} \leq \hat{E}_1 < \hat{E}_2 \leq 10 \text{ MeV}$ and $\hat{E}_2 < \hat{E}_3$ are satisfied. Note that the energy range of 10 keV to 10 MeV is an optimistic range of energies, and it is not known at the present time what energy range is allowed by current or future accelerator technology. **Figures 3, 4** display this same information but in the case of $K_p = 5$ and $K_p = 7$. From these calculations, we observe that $\hat{E}_3$ is significantly large on midnight local time, and T89 also predicts $B_\phi = 0$ on midnight. Therefore, we conclude from this that any assumptions that rely on Gardner's formula are valid on this domain. Additionally, we see that the region where this method may be applied is restricted to the case where the satellite is positioned near the equatorial plane. This may be problematic if the value of $\hat{E}_2$ is in the relativistic range, since μ may not necessarily be conserved for relativistic electrons in this region [see Willard et al. (2019) for further discussion on this topic]. However, as K_p increases, we see from **Figures 3, 4** that $\hat{E}_2$ decreases significantly. This indicates that the ability to infer field-line curvature may be possible from near the midplane during times where the field-line curvature is large. This is consistent with the dependence on R_c seen in the definition of $\hat{E}_2$ in (5b).

The precise situation where our method may be applied is as follows: a satellite equipped with an electron beam capable of firing at a range of energies will be in an orbit that intersects midnight local time near the midplane at a radial distance R and latitude θ_b . Once at midnight, the satellite will then begin to fire a beam with kinetic energy less than $\hat{E}_1$, and the precipitation of the beam in the ionosphere will be observed at a latitude θ_a . The time delay between the firing and the observation of the precipitation of the beam at the ionosphere is used to infer the total arc length of the field-line $\hat{\chi}$. The satellite will then gradually increase the energy of the beam until there is significantly reduced precipitation. We will assume that this critical energy is not greater than $\hat{E}_3$ so that the field-line curvature at the position of the satellite κ will also be inferred. We will also assume that all points on the field-line lie on midnight local time and that the magnetic field can be well-approximated as a dipole near the surface of the earth.

2. METHOD DESCRIPTION

2.1. Intuition

Our objective is to formalize a method by which field-lines may be approximated with curves that are consistent with measurements. In this section, we describe the problem that must be solved and the strategy that we take in solving it. The problem we would aim to solve is that of finding a family of curves that satisfy a set of constraints: (1) the curve must pass through the precipitation point and be consistent with a dipole field-line near the precipitation point, (2) the curve must pass through the position of the satellite and be consistent with the measured tangent direction and curvature of the field-line at the position of the satellite, and (3) the arc length of the curve between the launch and precipitation points must be consistent with the measured arc length. Our method satisfies these constraints by defining curves analytically near the end-points and defining the curves numerically over the rest of the domain. This strategy amounts to defining the curve piece-wise (see **Figure 5**), so that (a) the equation of the curve is exactly that of a dipole field-line near the precipitation point (satisfying the first constraint), (b) the equation of the curve is exactly a second-order polynomial near the launch point (satisfying the second constraint while assuming higher order derivatives are zero), and (c) the curve is numerically determined over the rest of the domain to satisfy the third constraint.

This strategy allows for the third constraint to be satisfied through a numerical procedure nearly independently of the first and second constraints. In this way, the problem is essentially simplified to the problem of finding families of curves having a given arc length between set end points. In order to solve this simplified problem, we let the curve be defined as an interpolation of a finite scatter of points. To understand how these points must be chosen to approximately satisfy the arc length constraint, it is best to consider a polygonal chain that has these points as vertices. By approximating the arc length along the curve between two of these prescribed points as merely the straight-line-distance between the points (see **Figure 6**), constraining the arc length of the curve in this sub-domain is approximately equivalent to constraining the total length of the chain. Our approximation method can then be understood to be that of an iterative process where the vertices of the chain are chosen one by one. At each step in the iteration, the choice of where the next vertex will be located is necessarily restricted, since at every stage of the iteration it is possible to choose a vertex that makes it impossible to finish constructing the chain without changing the length. This restriction can be clearly identified by considering the shortest possible chain connecting the two end points given a chosen vertex (see **Figure 7**). If this shortest possible chain has a total length greater than the required length, then there must not be a curve within the set of possible curves that passes through that chosen vertex. Restricting the vertex-choosing process in this way guarantees that, after carrying out this process through some number of iterations, the final vertex can always be chosen so that the required length of the chain may be satisfied exactly. The actual curve is then determined here

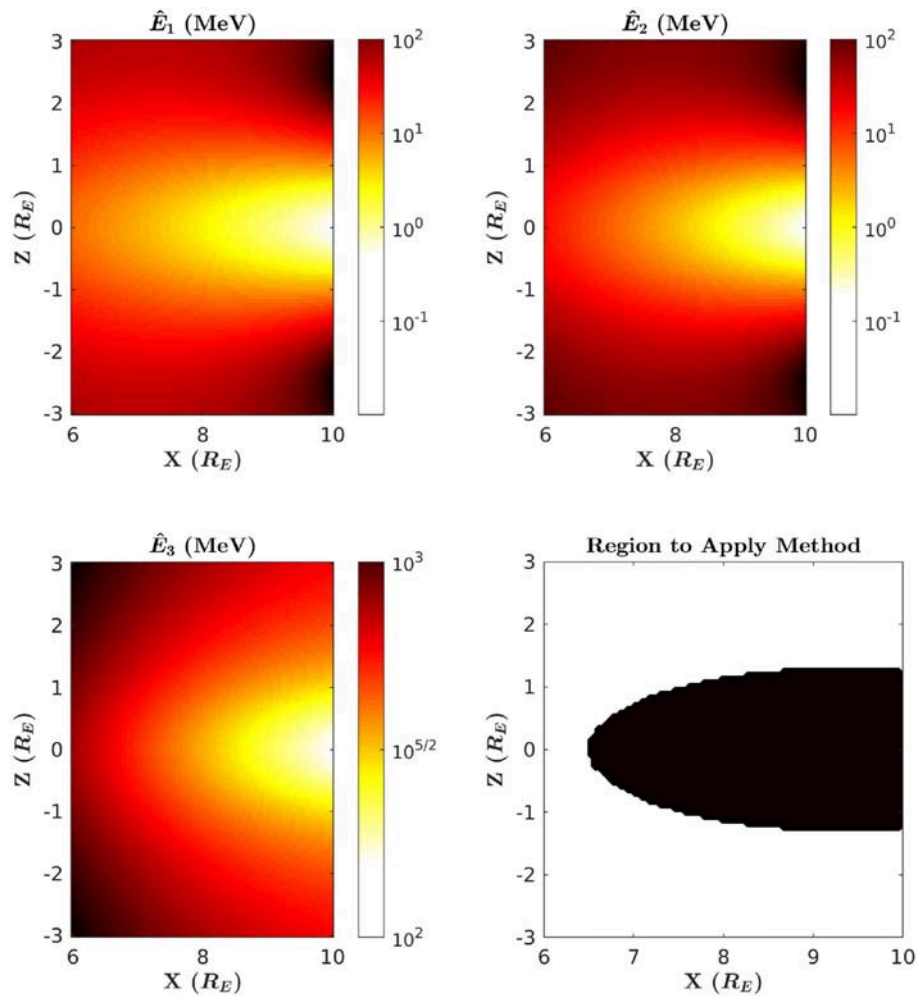


FIGURE 2 | $\hat{E}_1$, $\hat{E}_2$, and $\hat{E}_3$ computed over midnight local time in the case of $K_p = 1$. Bright colors indicate lower energy, and darker colors indicate higher energy. $\hat{E}_1$ is shown in the top left panel, $\hat{E}_2$ in the top right, and $\hat{E}_3$ in the bottom left. In the bottom right panel, the region where this method should be applied is shown in black.

to be an interpolation of the points generated by this iterative procedure. In this way of conceptualizing the method, the degrees of freedom seen in the general solution are manifested as the freedom to construct chains having any particular set of vertices so long as the choice of a particular vertex does not restrict the length of the chain connecting the end points to lengths greater than the required length.

To summarize, the problem of finding a general curve that is consistent with the information gathered by energy-variable electron beam experiments is not straightforward. With minimal loss in generality, we employ a strategy where curves are defined analytically near the end points and numerically over the rest of the domain. This allows for the arc length of the curve to be restricted nearly independently of the other constraints, which are localized about the end points. Our method for satisfying the arc length constraint is a numerical procedure where a finite set of points are each chosen iteratively, and the curve is ultimately given as an interpolation of these points.

2.2. Formalism

To represent the field-line, we define the function f such that all points (r, θ) on the field-line satisfy $r = f(\theta)$. The known or assumed geometric attributes of the field-line can then be expressed as constraints on f :

$$f(\theta_b) = R \quad (7a)$$

$$f'(\theta_b) = R \frac{B_r}{B_\theta} \quad (7b)$$

$$f''(\theta_b) = R \left(1 - R\kappa \left(1 + \frac{B_r^2}{B_\theta^2} \right)^{3/2} + 2 \frac{B_r^2}{B_\theta^2} \right) \quad (7c)$$

$$f(\theta) \approx R_E \frac{\sin^2(\theta)}{\sin^2(\theta_a)} \text{ for } \theta \approx \theta_a \quad (7d)$$

$$\int_{\theta_a}^{\theta_b} \sqrt{f^2 + f'^2} d\theta = \hat{\chi} \quad (7e)$$

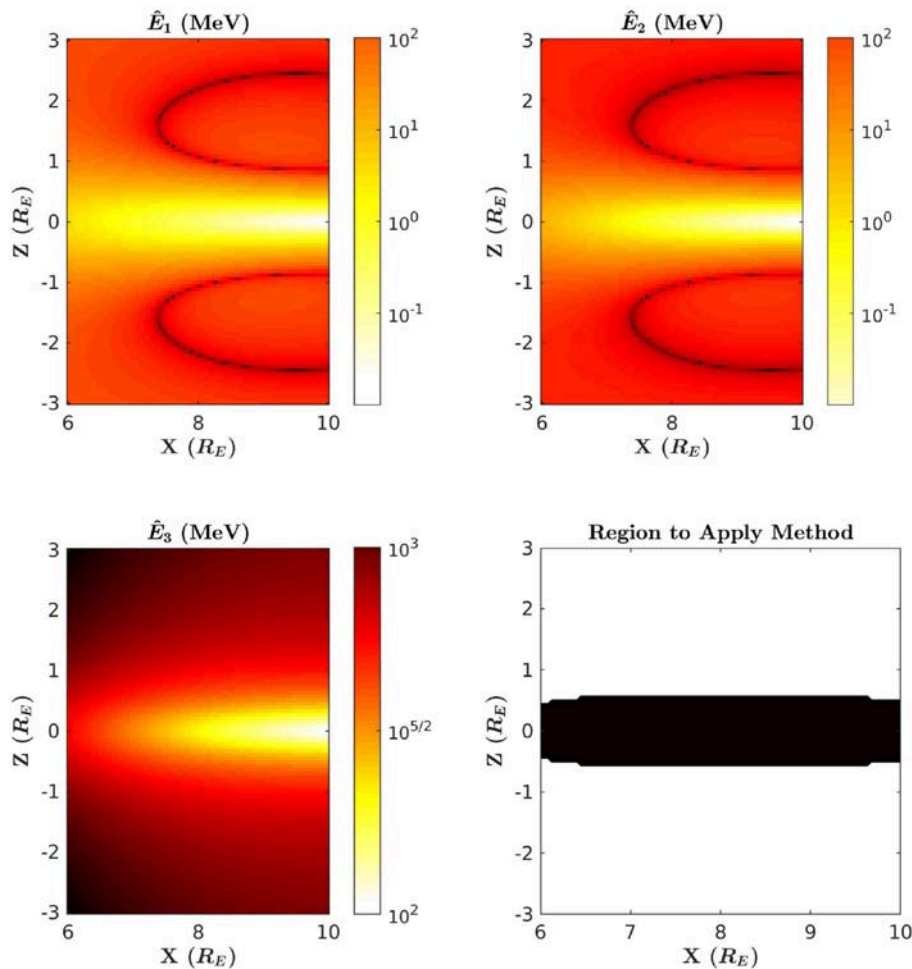


FIGURE 3 | $\hat{E}_1$, $\hat{E}_2$, and $\hat{E}_3$ computed over midnight local time in the case of $K_p = 5$. Bright colors indicate lower energy, and darker colors indicate higher energy. $\hat{E}_1$ is shown in the top left panel, $\hat{E}_2$ in the top right, and $\hat{E}_3$ in the bottom left. In the bottom right panel, the region where this method should be applied is shown in black.

Where R is the radial distance of the satellite, θ_a and θ_b are the latitude of the precipitation point and the position of the satellite, respectively, κ is the measured curvature, and $\hat{\chi}$ is the measured arc length. Note that the fourth constraint (7d) follows from the assumption that the field is well-approximated as a dipole near the Earth, since dipole field-lines are expressed as $r \propto \sin^2(\theta)$. These restrictions define the original problem mentioned previously. Following the strategy already described, we define f piece-wise:

$$f(\theta) = \begin{cases} f_1(\theta) & \text{for } \theta_a \leq \theta < \hat{\theta}_1 \\ f_2(\theta) & \text{for } \hat{\theta}_1 \leq \theta \leq \hat{\theta}_2 \\ f_3(\theta) & \text{for } \hat{\theta}_2 < \theta \leq \theta_b \end{cases} \quad (8)$$

where f_1 and f_3 are determined by only the first four of the above constraints (7a–7d), as well as the additional constraint that f_3

should have no higher order derivatives:

$$\begin{aligned} f_1(\theta) &= R_E \frac{\sin^2(\theta)}{\sin^2(\theta_a)} \\ f_3(\theta) &= R + (\theta - \theta_b) f'(\theta_b) + \frac{1}{2}(\theta - \theta_b)^2 f''(\theta_b) \end{aligned} \quad (9)$$

(see **Figure 5**). We define two length parameters L_1 and L_2 , which will allow the domain to be divided up with respect to distances:

$$\begin{aligned} \hat{\theta}_1 &= \sin^{-1}\left(\sqrt{\frac{L_1}{R_E}} \sin(\theta_a)\right) \\ \hat{\theta}_2 &= \theta_b + \sin^{-1}(L_2/R) \end{aligned} \quad (10)$$

L_1 and L_2 are distance scales representing how far from the origin the field-line can be expected to match f_1 and how far from the launch point it is expected to match f_3 , respectively. f_2 must

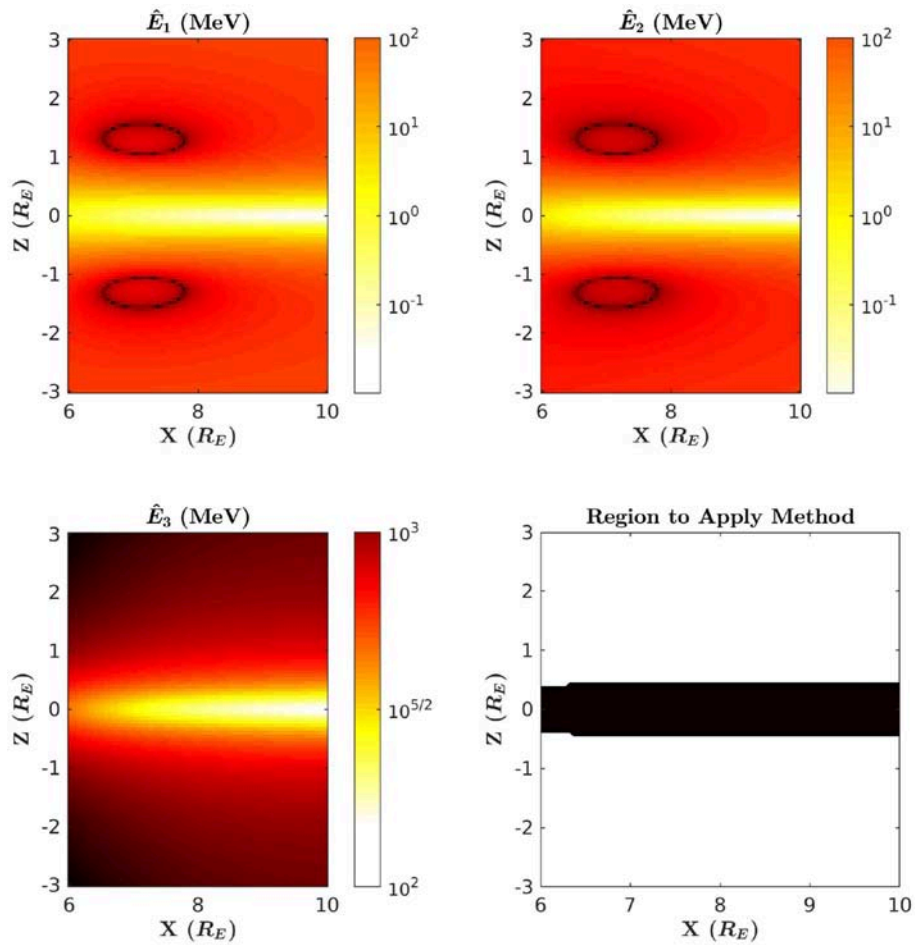


FIGURE 4 | $\hat{E}_1$, $\hat{E}_2$, and $\hat{E}_3$ computed over midnight local time in the case of $K_p = 7$. Bright colors indicate lower energy, and darker colors indicate higher energy. $\hat{E}_1$ is shown in the top left panel, $\hat{E}_2$ in the top right, and $\hat{E}_3$ in the bottom left. In the bottom right panel, the region where this method should be applied is shown in black.

satisfy:

$$f_2(\hat{\theta}_1) = f_1(\hat{\theta}_1), \quad f_2(\hat{\theta}_2) = f_3(\hat{\theta}_2) \quad (11)$$

$$\int_{\hat{\theta}_1}^{\hat{\theta}_2} \sqrt{f_2^2 + f_2'^2} d\theta = \tilde{\chi}$$

where we have defined a new arc length variable for brevity:

$$\tilde{\chi} = \hat{\chi} - \int_{\theta_a}^{\hat{\theta}_1} \sqrt{f_1^2 + f_1'^2} d\theta - \int_{\hat{\theta}_2}^{\theta_b} \sqrt{f_3^2 + f_3'^2} d\theta \quad (12)$$

We define L_f to be a function giving the straight line distance between two points on the curve $r = f(\theta)$:

$$L_f(a, b) = \left\| \begin{pmatrix} f(b) \sin(b) \\ f(b) \cos(b) \end{pmatrix} - \begin{pmatrix} f(a) \sin(a) \\ f(a) \cos(a) \end{pmatrix} \right\| \quad (13)$$

for arbitrary angles a and b (see **Figure 6**). If we consider a finite set of angles $\{\Theta_1, \dots, \Theta_N\}$ that are evenly spaced over the domain

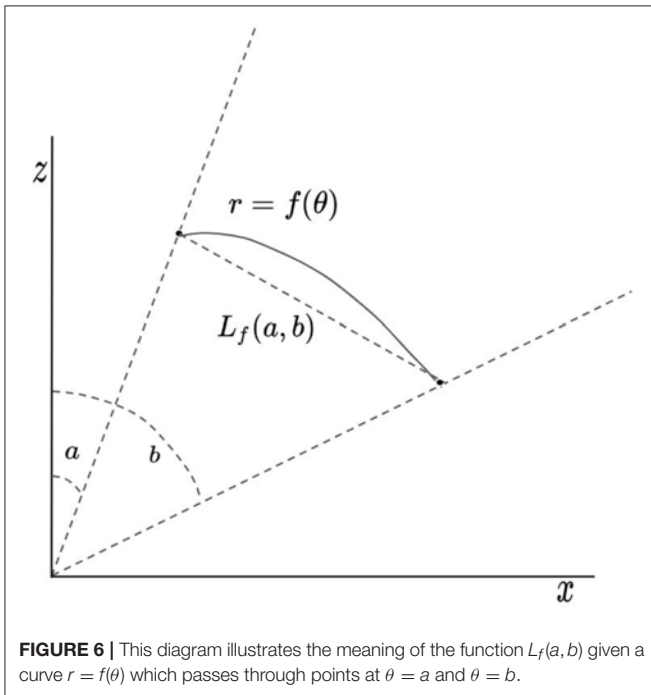
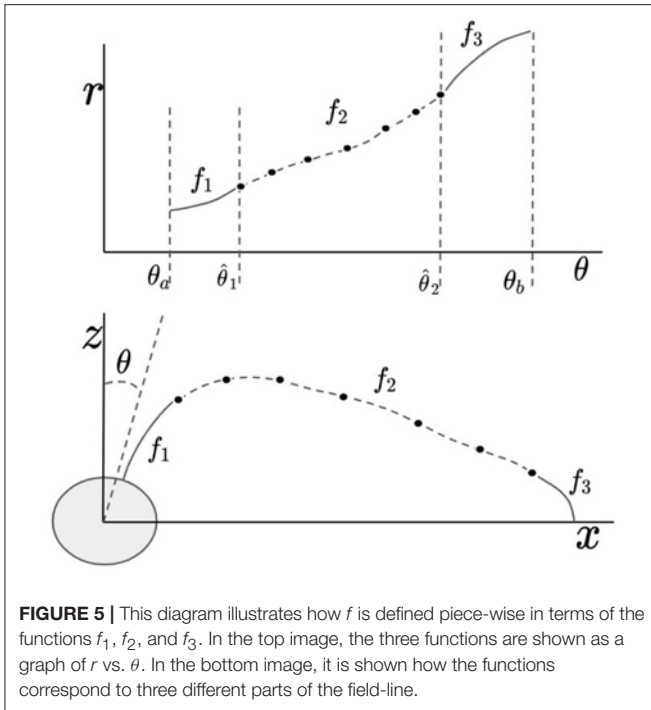
$[\hat{\theta}_1, \hat{\theta}_2]$, and N is chosen such that the discretization is sufficiently fine, then the integral equation may be well-approximated by

$$\sum_{i=1}^{N-1} L_f(\Theta_i, \Theta_{i+1}) = \tilde{\chi} \quad (14)$$

At each angle Θ_n , we define a function $D_n(f)$ by

$$D_n(f) = \sum_{i=1}^n L_f(\Theta_i, \Theta_{i+1}) + L_f(\Theta_{n+1}, \Theta_N) \quad (15)$$

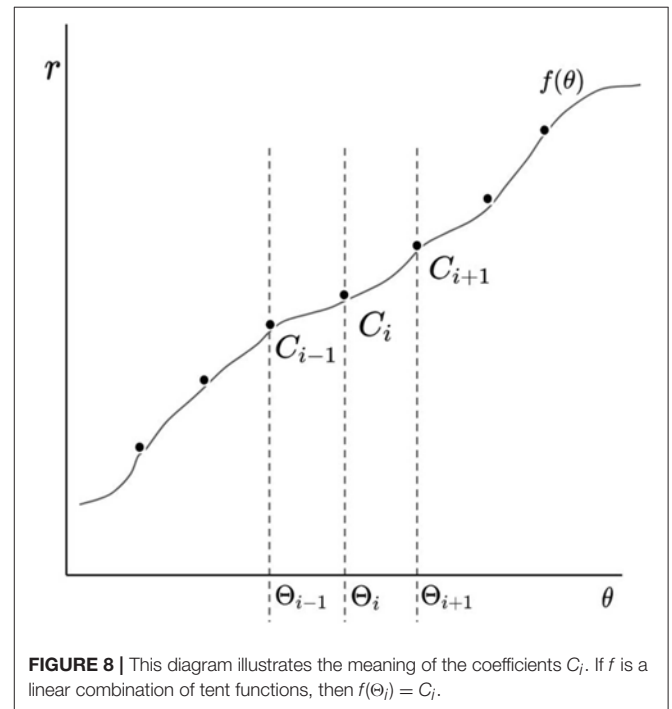
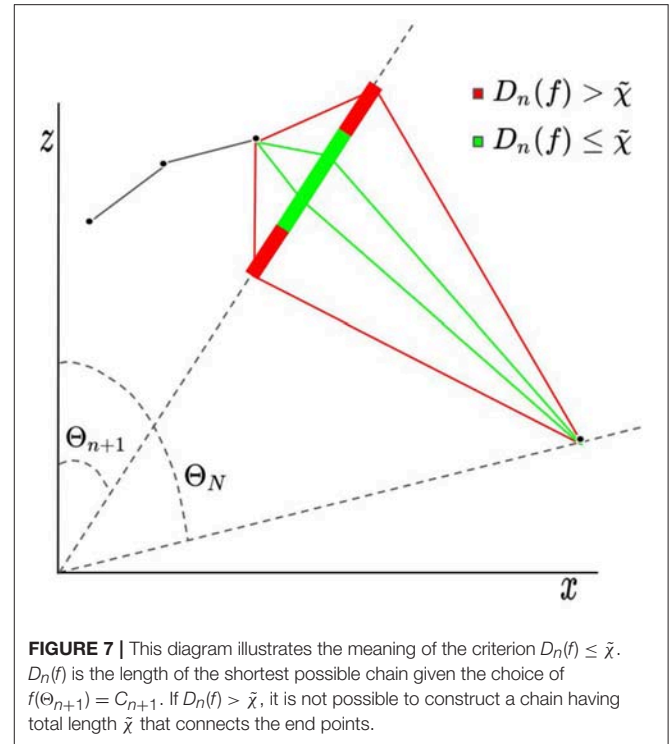
$D_n(f)$ is the length of the chain having vertices at each of the angles $\{\Theta_1, \dots, \Theta_{n+1}, \Theta_N\}$. $D_n(f)$ represents the minimum length of a chain given a chosen vertex at the angle Θ_{n+1} . In order for the chosen vertex to be allowed, one can check that $D_n(f) \leq \tilde{\chi}$ (see **Figure 7**). We restrict f_2 to be a linear combination of



functions ϕ_i , which are commonly known as tent functions:

$$f_2(\theta) = \sum_{i=1}^N C_i \phi_i(\theta) \quad (16)$$

$$\phi_i(\theta) = \frac{N}{\hat{\theta}_2 - \hat{\theta}_1} \begin{cases} \Theta_{i+1} - \theta & \text{for } \Theta_i \leq \theta < \Theta_{i+1} \\ \theta - \Theta_{i-1} & \text{for } \Theta_{i-1} \leq \theta < \Theta_i \\ 0 & \text{otherwise} \end{cases}$$



Note that $f_2(\Theta_i) = C_i$ (see **Figure 8**). Given the reasoning previously described in terms of polygonal chains, the coefficients may be chosen so that they satisfy the following recursion rule:

$$C_1 = f_1(\hat{\theta}_1) \quad (17a)$$

$$C_{n+1} \in \left\{ c \mid D_n \left(\sum_{i=1}^n C_i \phi_i + c \phi_{n+1} \right) \leq \tilde{\chi} \right\} \quad (17b)$$

$$C_{N-1} \in \left\{ c \mid D_{N-2} \left(\sum_{i=1}^{N-2} C_i \phi_i + c \phi_{N-1} \right) \leq \tilde{\chi} \right\} \quad (17c)$$

$$C_N = f_3(\hat{\theta}_2) \quad (17d)$$

The degrees of freedom in the general solution are here expressed as the freedom to choose any set of coefficients that satisfy the above relations. This freedom must now be expressed in terms of some finite set of parameters. There are undoubtedly many possible approaches that one could take in doing this. To prove that this is possible, one can check that this can be done simply by replacing the above recursion rule (17b) with the formula:

$$C_{n+1} = \min(Z_n) + \frac{1}{2} \left(\max(Z_n) - \min(Z_n) \right) (\tanh(K_n) + 1)$$

$$Z_n = \left\{ c \mid D_n \left(\sum_{i=1}^n C_i \phi_i + c \phi_{n+1} \right) \leq \tilde{\chi} \right\} \quad (18)$$

Any chosen set of real numbers $\{K_1, \dots, K_{N-3}\}$ correspond to a particular solution to the problem.

3. EXAMPLE APPLICATIONS

In this section, we will show the approximations that are generated from our method when the parameters $\hat{\chi}$ and κ are taken from realistic field-lines obtained from T89 on midnight local time with parameters chosen to correspond with $K_p = 1$. Although the method is consistent, it must be shown that realistic curves can be easily obtained by imposing realistic restrictions on the remaining degrees of freedom. For this purpose, it is not necessary to express these degrees of freedom in terms of any free parameters as described in the previous section. Rather, it is sufficient for our purposes to include an additional restriction to the recursion:

$$C_{n+1} \in \left\{ C_n \left(1 + \left(2 \frac{l}{M} - 1 \right) \epsilon \right) \right\} \quad (19)$$

Where $l = 0, 1, \dots, M$ and ϵ sets a maximum fractional increase between C_n and C_{n+1} . The set of all particular solutions to the problem is now guaranteed to be a finite set of functions. Through a brute force algorithm, we may then systematically generate each particular solution and then sort them by the average square second derivative of f_2 , which is equivalent to sorting by

$$\langle f_2'' \rangle \sim \sum_{i=2}^{N-1} (C_{i+1} + C_{i-1} - 2C_i)^2 \quad (20)$$

Field-lines with small f'' are typical of T89, so it is expected that choosing f_2 as the curve with the least $\langle f_2'' \rangle$ should result in curves that are not very different from the original field-lines from which

the parameters $\hat{\chi}$ and κ were obtained. Field-lines can be traced from T89 as parametric curves $(x(s), z(s))$ satisfying:

$$\begin{pmatrix} x'(s) \\ z'(s) \end{pmatrix} = \hat{b}(x(s), z(s)) \quad (21)$$

where $\hat{b} = \vec{B}/B$. Here, s represents the arc length between the point $(x(s), z(s))$ and the point $(x(0), z(0))$. The field-lines are traced by iteratively solving the above equation from a chosen launch position until $\|(x(s_f), z(s_f))\| \leq 1$, for some value s_f , at which point it is clear that $\hat{\chi} \approx s_f$. The curvature at the launch point is then computed from the formula $\kappa = (\hat{b} \cdot \nabla) \hat{b}$.

For this demonstration, we will make the choices $L_1 = 3R_E$, $N = 10$, $M = 15$, and $\epsilon = 0.2$. Our choice of L_2 is different depending on the launch point: for launch points at the midplane, we choose $L_2 = R_c$, while off the midplane the value is chosen more conservatively to be $L_2 = R_E/3$. **Figure 9** shows four examples of the smoothest solutions generated with this method alongside the T89 field-lines used to obtain the parameters. **Figure 10** shows the ten smoothest generated curves only in comparison with each other so that the remaining degrees of freedom can be visualized. **Figure 11** shows the same information as **Figures 9, 10**, but is an example of using launch points that are not on the equatorial plane and are instead slightly away from the equatorial plane. These examples show that this method can be easily constrained to produce realistic field-lines that match well with the original field-lines used to obtain the parameters.

4. DISCUSSION

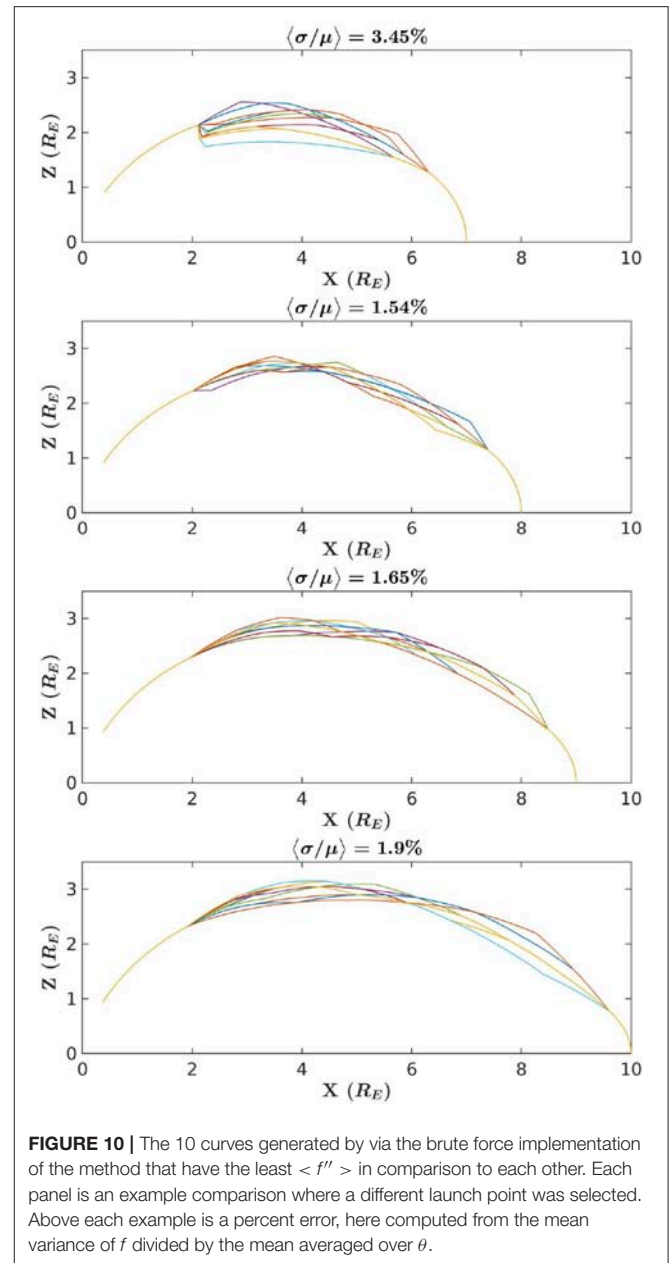
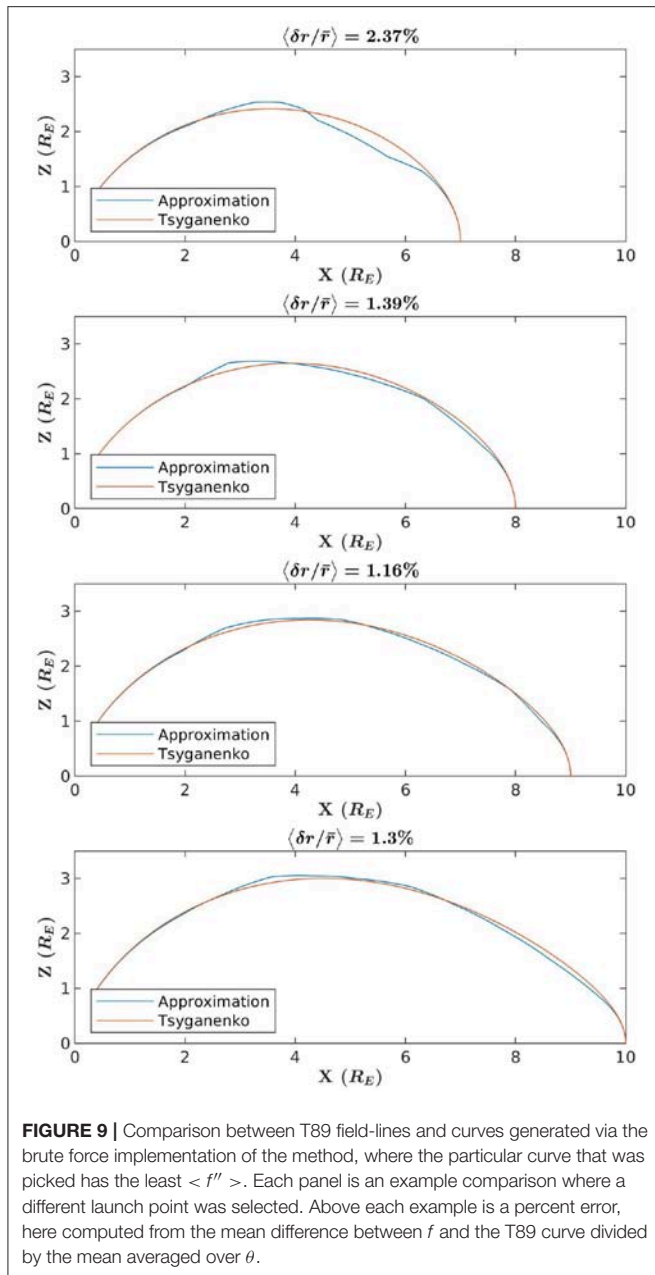
In the above example, we show that the degrees of freedom seen in the family of curves may be easily constrained to agree well with T89 by minimizing the average square of the second derivative of f . In actual practice, this method should be used as part of the three step data assimilation technique mentioned previously. As an example, suppose the model that we wish to constrain is the axisymmetric Grad-Shafranov equation:

$$\nabla \cdot \left(\frac{1}{r^2 \sin^2(\theta)} \nabla \psi \right) = -\mu_0 \frac{dP}{d\psi} \quad (22)$$

where the magnetic field is $\vec{B} = (\nabla \psi \times \hat{\phi})/r \sin(\theta)$, P is the plasma pressure, and we have assumed $B_\phi = 0$. The first step of the procedure would then be to find a solution to Grad-Shafranov that optimizes some set of measurements to obtain a naive forecast potential ψ . In the second step, ψ would be assimilated to optimize the field-line approximation by minimizing a cost function J related to the variance of ψ on the curve $r = f(\theta)$:

$$J \sim \langle w(\theta) \psi(\theta, f(\theta)) - \langle \psi(\theta, f(\theta)) \rangle \rangle \quad (23)$$

where $w(\theta)$ is some weight function. The final step of the three step data assimilation approach would be to then assimilate the curve $r = f(\theta)$ together with the equilibrium model. This may be done by finding an optimal solution as in the first step



only with the added constraint that $\psi(f(\theta), \theta) = \text{constant}$. This example is particularly elegant, since the added constraint is a Dirichlet boundary condition, provided that the domain of the calculation is restricted to the region enclosed by the approximated field-line.

Further investigation is necessary in order to fully justify the experimental techniques described at the start of this paper. The ability to infer the curvature of field-lines relies on the ability to accurately aim the electron beam, and it has yet to be determined how feasible this is given the current technology. As mentioned previously, this inference also relies on the assumption that Gardner's asymptotic expansion of μ is conserved, and the degree to which this assumption

can still be made given perturbations of the magnetic field has yet to be determined. It is also a possibility that artificially injecting electrons into the ambient plasma may drive instabilities that will significantly affect the path of the beam. Although past experiments have shown that this possibility is not necessarily significant (Winckler, 1980), further investigation is necessary in order to fully determine which conditions would require that the ambient plasma is taken into account.

The task of designing energy-variable electron beam experiments certainly has many difficult challenges that must be overcome, so it is necessary to consider the importance of this feature with respect to our method. Primarily, the

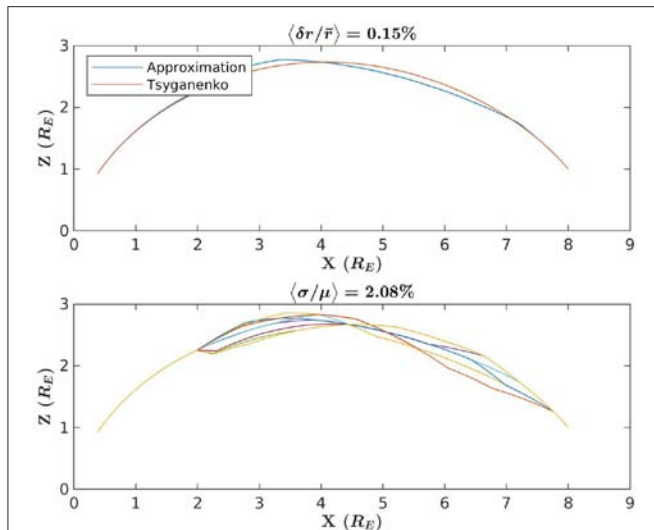


FIGURE 11 | The top image conveys the same information as the images in **Figure 9**, but the launch point for this example is chosen to be away from the equatorial plane. The bottom image likewise conveys the same information as the images in **Figure 10** but for a launch point away from the equatorial plane.

ability to vary the energy of the beam is required in order to directly measure the field-line curvature at the launch position. However, curvature is also necessary in order to know the value of $\hat{E}_1$, and therefore a lack of knowledge of the curvature leads to a source of uncertainty in the arc length measurement. Without an energy-variable experiment, it would therefore be necessary to simply infer the field-line curvature by some other means. For example, if the satellite is positioned well above the equatorial plane, it may be reasonable to simply assume that $\kappa \approx 0$ and $\hat{E}_1$ is relatively large. However, it is uncertain as to whether or not this application would be of much use to modeling the inner magnetosphere if the field-line curvature is not actually measured.

As mentioned previously, this method as a whole is only applicable in those cases where the arc length and curvature may be measured using a single satellite equipped with an energy-variable accelerator. However, various techniques and concepts employed in this method may be adapted to be used in alternative cases. For example, the technique used to approximate the arc length constraint using polygonal chains may be adapted to any context where the arc length of field-lines is known. Provided some method by which the field-line torsion may be inferred, we may additionally consider generalizing this method to study field-lines that are not restricted to midnight local time. This method should therefore be seen as a particular implementation of a more general approach to

utilizing measurements of field-line geometry that may utilize a wider variety of measurements than discussed here.

5. SUMMARY

In this article, we discuss a way in which it may be possible to measure field-line curvature and arc length using energy-variable electron beam experiments. In order to use these measurements to constrain models of Earth's inner magnetosphere, we discuss a three step data assimilation approach where prior information about the field may be used to approximate the field-line. The prior information may then be used in conjunction with this approximation to better constrain the model. Central to this approach is the method presented in this article, which is a means of approximating a general solution to a set of constraints, such that the problem is only slightly more restricted than mere adherence to the measurements of field-line arc length and curvature, and the degrees of freedom in this general solution can be expressed in terms of free variables. As an example, we obtain parameters from realistic field-lines traced from T89, and compare these curves with select approximations generated using the method to show that T89 curves can be reproduced to good accuracy using the method by imposing a realistic bias.

DATA AVAILABILITY

The raw data supporting the conclusions of this manuscript will be made available by the authors, without undue reservation, to any qualified researcher.

AUTHOR CONTRIBUTIONS

JW wrote the manuscript and created necessary programming tools with the help from JJ and JS. JS devised the chain approximation approach to satisfying arc length constraints. JW and JJ devised the data assimilation procedure. JW developed the formalism in consultation with JJ, JS, AP, IK, and ES. JW, JJ, JS, AP, IK, and ES contributed to discussing the scope and advantages of the method.

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Evolution of a Relativistic Electron Beam for Tracing Magnetospheric Field Lines

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Tracing magnetic field-lines of the Earth's magnetosphere using beams of relativistic electrons will open up new insights into space weather and magnetospheric physics. Analytic models and a single-particle-motion code were used to explore the dynamics of an electron beam emitted from an orbiting satellite and propagating until impact with the Earth. The impact location of the beam on the upper atmosphere is strongly influenced by magnetospheric conditions, shifting up to several degrees in latitude between different phases of a simulated storm. The beam density cross-section evolves due to cyclotron motion of the beam centroid and oscillations of the beam envelope. The impact density profile is ring shaped, with major radius ~ 22 m, given by the final cyclotron radius of the beam centroid, and ring thickness ~ 2 m given by the final beam envelope. Motion of the satellite may also act to spread the beam, however it will remain sufficiently focused for detection by ground-based optical and radio detectors. An array of such ground stations will be able to detect shifts in impact location of the beam, and thereby infer information regarding magnetospheric conditions.

Keywords: relativistic particle beam, beam envelope, nonneutral plasmas, electron beams (e-beams), field-line mapping, computational modeling, ballistic simulation, active space experiments

1. INTRODUCTION

The injection of artificial electron beams into the Earth's magnetosphere has proven to be a powerful diagnostic tool for studying the physics of the magnetosphere, ionosphere and upper atmosphere (Winckler, 1980). A large number of experiments have focused on the near plasma environment of the ionosphere and chemistry of the upper atmosphere, however three (known) experiments have injected beams from sounding rockets upwards into the magnetosphere. The *Hess Artificial Aurora Experiments* (Hess et al., 1971) and the joint French-Soviet *ARAKS Experiment* (Gendrin, 1974) observed atmospheric emission on the opposite hemisphere to where particles were injected, indicating that electron beams could survive a transition through the magnetosphere. The *ECHO* experiments (Hendrickson et al., 1971) utilized detectors near to the injection location, demonstrating that particles could undergo multiple transitions from hemisphere to hemisphere, maintaining beam stability and detectability. In all, seven *ECHO* experiments were performed, providing unique insight into the workings of the magnetosphere (Winckler, 1982; Winckler et al., 1989).

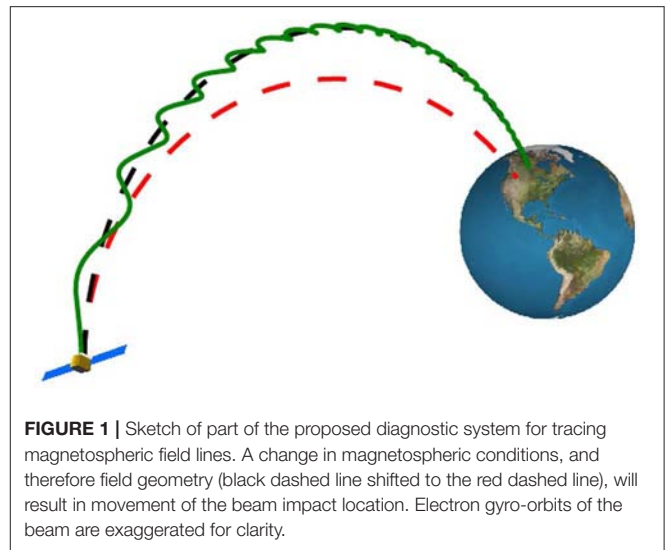
These earlier experiments injected beams with energies < 40 keV, however advances in accelerator technology now make it feasible for spacecraft to generate beams of electrons with relativistic energies > 0.5 MeV (Banks et al., 1987; Mishin, 2005). For a fixed beam current, relativistic beams result in reduced spacecraft charging due to lower beam density requirements. Furthermore, three-dimensional particle-in-cell simulations have shown that relativistic beams are more stable than lower energy beams during emission from a spacecraft (Gilchrist et al., 2001; Neubert and Gilchrist, 2002, 2004). It has been proposed that relativistic electron beams could be an ideal diagnostic for field-line tracing within the magnetosphere, assisting in the validation and development of advanced magnetospheric models. Such a diagnostic may also provide additional insights via the active modification of the space-plasma environment (National Research Council, 2013).

The advent of high-power, low-voltage RF amplifier chips, such as high-electron-mobility transistors, has enabled the development of new electron linear accelerator technologies (Lewellen et al., 2018). Each accelerator cavity can be coupled to its own lightweight, compact amplifier as opposed to the entire device being powered by a heavier high-voltage klystron (Nguyen et al., 2018), resulting in a comparatively lighter and more robust device. The lower mass and power requirements make it feasible to mount such an accelerator onto a spaceborne satellite. Efforts being undertaken at the Los Alamos National Laboratory, SLAC National Accelerator Laboratory and Goddard Space Flight center have worked to characterize the RF amplifier performance, optimize the accelerator structure, demonstrate radiation hardness and conduct an experimental technology validation program (Lewellen et al., 2018; Nguyen et al., 2018).

Attached to an orbiting satellite, such a compact linear accelerator could launch relativistic electrons onto various field lines of the magnetosphere over a range of magnetospheric conditions. In an ideal scenario, electrons launched from the satellite will trace the field-lines of the magnetosphere until precipitation in the upper atmosphere. Precipitating electrons then produce optical emission and density enhancement signatures in the D-region of the atmosphere, detectable by an array of ground stations (Marshall et al., 2014). Therefore the diagnostic system consists of an orbiting satellite, compact linear accelerator, and numerous ground stations, likely coordinated by a central control system. A sketch of the satellite, electron beam, and impact location over the North American continent is shown in **Figure 1**.

In regards to beam propagation following injection into the magnetosphere and until precipitation in the upper atmosphere, there are three primary questions:

1. Will the injected electrons reach the upper atmosphere?
2. Will the changing magnetospheric conditions influence the impact location of electrons such that this change is detectable by ground stations?
3. Will the beam density profile at impact with the upper atmosphere be sufficiently narrow to produce an emission signature distinguishable from background noise?



With respect to Question 1, the most fundamental consideration is whether the magnetic field line along which a particle is injected will intersect the Earth. For an arbitrary injection location and unknown field geometry, this is difficult to determine a priori. However, close to the Earth ($< 10R_E$, where R_E is the radius of the Earth) the field geometry is close to that of a dipole, and therefore particles launched from near the geomagnetic equator will most likely be attached to field lines which intersect the Earth. A particle injected onto such a field line will experience an increasing magnetic field strength as it approaches the Earth. If the particle has an initially non-zero magnetic moment, then conservation of magnetic moment and energy will result in parallel kinetic energy being converted into perpendicular kinetic energy during this transition. If the increase in field strength is sufficient, then all of the initial energy may be converted to perpendicular energy and the particle is mirrored at a location outside of the Earth's atmosphere, thus precluding precipitation into the atmosphere. In the language of a magnetic mirror, we consider particles which are initialized such that their mirror radius is smaller than the radius of the Earth to be within the loss cone.

A perfectly uniform, non-relativistic beam injected directly along the field line will have zero magnetic moment, and will therefore always precipitate. A realistic beam, however, always has a finite perpendicular energy spread, known as beam *emittance* within the literature (see section 2.3 for a formal definition), thus particles will have a non-zero magnetic moment. For relativistic electrons, it is important to consider higher order components of the magnetic moment asymptotic expansion when determining the mirror point. Porazik et al. (2014) shows that the loss cone of viable injection angles narrows for increasing beam energy. Furthermore, for injection from the equatorial plane within the magnetotail, the loss cone is narrowed with increased dipole stretching (and therefore local curvature). In such a geometry it becomes favorable to inject above the equatorial plane where field lines have reduced curvature (Willard et al., 2019).

These findings set limitations on the beam energy, injection location, pointing precision and beam emittance for a viable diagnostic system.

This paper seeks to provide insight into Questions 2 & 3 via numerical and analytic analysis of electron beam propagation. The following section discusses the methodology of our analysis and numerical tools. Sections 3, 4 present results which pertain to Questions 2 and 3, respectively. For further details and a more complete picture of this proposed diagnostic, see Sanchez et al. (in preparation).

2. METHODOLOGY

2.1. Coordinate Systems

Due to the vastly different length scales and the symmetries of a beam and the magnetosphere, it is necessary to depend on four different coordinate systems. **Figure 2** shows how each of these systems are interrelated. The first is a Cartesian coordinate system with basis vectors $\{\hat{\mathbf{X}}, \hat{\mathbf{Y}}, \hat{\mathbf{Z}}\}$ and origin at the center of the Earth. The $\hat{\mathbf{X}}$ direction points from the center of the Earth toward the center of the Sun. The $\hat{\mathbf{Z}}$ direction points along geomagnetic North, and $\hat{\mathbf{Y}} = \hat{\mathbf{Z}} \times \hat{\mathbf{X}}$. This is generally referred to as the solar magnetic or centered dipole coordinate system.

The second coordinate system is spherical with basis vectors $\{\hat{\mathbf{R}}, \hat{\theta}, \hat{\phi}\}$ and origin collocated with that of the $\{\hat{\mathbf{X}}, \hat{\mathbf{Y}}, \hat{\mathbf{Z}}\}$ system. $\hat{\mathbf{R}}$ points in the radial direction, $\hat{\theta}$ is the polar angle, measured from the $\hat{\mathbf{Z}}$ axis and $\hat{\phi}$ is the azimuthal angle measured from the $\hat{\mathbf{X}}$ axis. At the Earth's surface, the angles $90^\circ - \theta^\circ$ and ϕ° correspond to angles of latitude and longitude, respectively, in geomagnetic coordinates.

The third coordinate system is Cartesian with basis vectors $\{\hat{\mathbf{x}}, \hat{\mathbf{y}}, \hat{\mathbf{z}}\}$ and origin at the location of a moving test particle, or average location of a collection of particles which make up a beam. The $\hat{\mathbf{z}}$ vector points along the local magnetic field line, the $\hat{\mathbf{x}}$ vector is normal to $\hat{\mathbf{z}}$ and lies along the $\hat{\mathbf{X}} \times \hat{\mathbf{Y}}$ plane, and $\hat{\mathbf{y}} = \hat{\mathbf{z}} \times \hat{\mathbf{x}}$.

The fourth coordinate system is spherical with basis vectors $\{\hat{\mathbf{v}}, \hat{\delta}, \hat{\lambda}\}$ and describes the particle velocity vector with respect to the local magnetic field. The origin of the system is collocated with the $\{\hat{\mathbf{x}}, \hat{\mathbf{y}}, \hat{\mathbf{z}}\}$ system. $\hat{\delta}$ is the polar angle measured from the $\hat{\mathbf{z}}$ axis and $\hat{\lambda}$ is the azimuthal angle measured with respect to the $-\hat{\mathbf{y}}$ direction.

2.2. Magnetic Field Geometry

For much of this theoretical analysis, a simple dipole model for the Earth's magnetosphere is used, centered in the Cartesian or spherical Earth-based coordinate systems and with dipole moment $\mathbf{D}$ pointing along the $-\hat{\mathbf{Z}}$ axis. This field has magnetic vector potential,

$$\mathbf{A} = -\frac{D}{R^2} \sin \theta \hat{\mathbf{Z}} \quad (1)$$

where, for a best fit with the Earth's magnetosphere $D = |\mathbf{D}| = 8.60 \times 10^{15} \text{ T} \cdot \text{m}^3$ (Porazik et al., 2014). This yields magnetic field components and magnitude,

$$\mathbf{B} = \nabla \times \mathbf{A} = -\frac{D}{R^3} (2 \cos \theta \hat{\mathbf{X}} + \sin \theta \hat{\mathbf{Y}}) \quad (2a)$$

$$B = |\mathbf{B}| = D \frac{\sqrt{4 - 3 \sin^2 \theta}}{R^3}, \quad (2b)$$

which describe field lines with profile,

$$R = R_0 \sin^2 \theta \quad (3)$$

where R_0 is the point at which the field line intersects the magnetic equatorial plane (the $\hat{\mathbf{X}} \times \hat{\mathbf{Y}}$ plane), also commonly known as L within the literature. The polar angle with which a field line intersects the Earth's surface θ_E can be found by inverting Equation (3) with $R = R_E$, where $R_E = 6,371 \text{ km}$ is the radius of the Earth,

$$\theta_E = \sin^{-1} \sqrt{\frac{R_E}{R_0}} \quad (4)$$

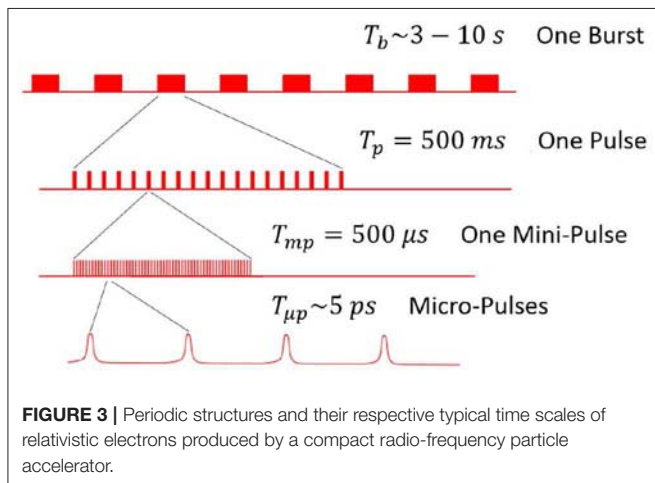
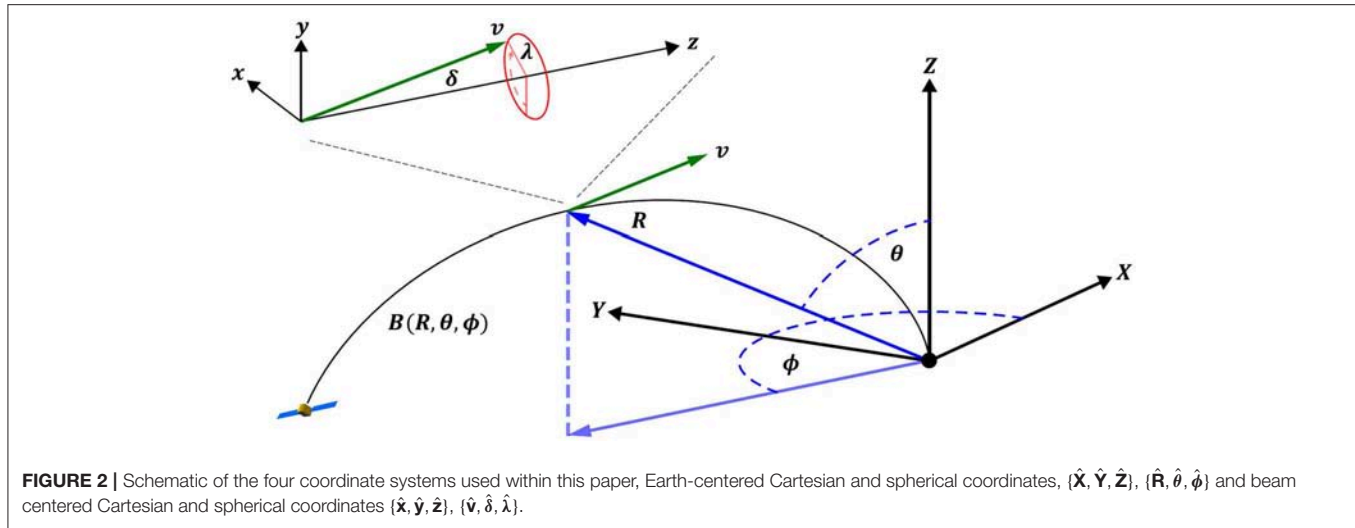
In addition to this simple model, realistic semi-empirical magnetic field geometry is implemented to study the effect of different magnetospheric conditions on particle trajectories. This geometry is implemented via the BATS-R-US (Powell et al., 1999; Tóth et al., 2012) package, as part of the Space Weather Modeling Framework (Tóth et al., 2005). The package solves for field geometry via the three-dimensional magneto-hydrodynamic equations on an adaptive grid with solar wind input data being supplied by the NASA Advanced Composition Explorer satellite (Stone et al., 1998).

2.3. Beam Parameters

A compact satellite-mounted particle accelerator of the radio-frequency (RF) type produces an electron beam which is non-uniform along the direction of propagation. The beam consists of periodic structures with various time and length scales (see **Figure 3**). The smallest scale structures are so called “micro-pulses” which are synchronized with the RF cycle. Multiple micro-pulses form a “mini-pulse,” a collection of which then forms one “pulse.” The timing of each mini-pulse, pulse and the subsequent entire “burst” of each beam firing is determined by the physics of the particle accelerator, spacecraft power limitations and scientific goals of the mission.

In this paper we consider down to the time scale of a mini-pulse since for typical energy spreads of RF accelerators the micro-pulse structure will quickly become indistinguishable over the path lengths considered. The largest time scale considered is that of a single pulse, since mission requirements demand that a single pulse impacting the atmosphere be detectable by ground stations.

The electron beam reference conditions of **Table 1** are used throughout much of this work.

**TABLE 1** | Physical parameters of the reference relativistic electron beam.

Property	Symbol	Value	Units
Beam energy	E_0	1.0	MeV
Beam beta	β	0.941	—
Beam lorentz factor	γ	2.957	—
Beam energy spread	ΔE	0.01	—
Radial emittance	ε_r	1.0	mm — mrad
Average beam current	I_0	1	mA
Initial mini-pulse time	T_{mp}	500	μs
Initial mini-pulse length	L_{mp}	141.1	km
Mini-pulse duty cycle	τ_{dc}	0.1	—
Initial beam radius	$r_{b,i}$	2	mm
Initial pulse time	T_p	0.5	s

Based on these properties, the initial phase space profile of a mini-pulse aligned with the local magnetic field vector is given by,

$$f(x, y, z, v_x, v_y, v_z, t = 0) = f_0 \exp \left(-\frac{x^2 + y^2}{2r_{b,i}^2} - \frac{z^2}{2(L_{mp}/2)^2} \right) \exp \left(-\frac{v_x^2 + v_y^2}{2\langle v_{\perp} \rangle^2} - \frac{(v_z - v_0)^2}{2\langle v_{\parallel} \rangle^2} \right), \quad (5)$$

where $v_0 = \beta c$, $\langle v_{\parallel} \rangle = q_e E_0 \Delta E / m_e c \beta \gamma^3$ is the RMS longitudinal velocity spread, and $\langle v_{\perp} \rangle = \varepsilon_r v_0 / r_{b,i}$ is the RMS radial velocity spread. Here ε_r is the beam emittance (note that the symbol ε_r is used to distinguish it from the small parameter ϵ used throughout this text). For a beam not initially aligned with the local field, the profile of Equation (5) will be tilted via angles δ and λ as per Figure 2.

2.4. Numerical Methods

For the numerical aspect of this work, simulations were performed with a single-particle-motion, ballistic propagation code, first used and verified in Porazik et al. (2014). The code initializes one beam pulse, consisting of a fixed number of mini-pulses. Each mini-pulse consists of a statistical distribution of particles spread in six-dimensional phase space (see Equation 5) and initialized via a pseudo-random number generator. The particles are evolved via the standard Boris algorithm (Boris, 1970) from some injection location (X, Y, Z) with provided injection angle (δ, λ) , along a prescribed magnetic field, until each particle has impacted the Earth ($R < R_E$) or the simulation time ends.

The code can incorporate either an analytic magnetic dipole field or take BATS-R-US data as an input. In the case of BATS-R-US data, magnetic field line information is interpolated to the particle location via a three-dimensional spline interpolation tool.

A limitation of the ballistic code is that it does not capture self-consistent collective interactions between the electron beam

and itself, or the ambient plasma. In section 4.4, we demonstrate that for beam properties near those of the reference conditions (see **Table 1**), the interaction of the beam with itself plays a small role compared to the independent motion of each particle. Interaction between the beam and the ambient plasma, however, may result in instabilities which act to spread the beam. Simple linear analysis suggests that a beam propagating through the magnetosphere will be stable to two-stream instabilities (Galvez and Borovsky, 1988), resistive hose, ion hose and filamentation instabilities (Gilchrist et al., 2001). A more detailed non-linear analysis is required and will be reserved for a future publication.

3. BEAM IMPACT LOCATION

In this section we consider the second question posed in the introduction to this paper, whether changing magnetospheric conditions will influence the impact location of the beam such that this change will be detectable by ground stations.

The section begins by studying the motion of a single electron with energy E_0 injected onto a dipole field line until it impacts with the Earth. We show that it is reasonable to assume conservation of the magnetic moment, an important result for the theoretical results of this paper. We then measure the offset of the final impact location with respect to the original field line due to particle drifts and compare this to analytic theory. Realistic magnetic field geometry is then considered, demonstrating that changes in magnetospheric conditions will appreciably shift the beam impact location.

3.1. Effect of Single-Particle-Motion Drifts

A single electron is injected from $-10R_E\hat{\mathbf{x}}$ with energy $E_0 = 1$ MeV from **Table 1** and an initial velocity vector along the local dipole field line (in the $\hat{\mathbf{z}}$ direction). The initial field line and trajectory of the electron from injection until impact with the Earth is shown on the three Cartesian planes in **Figure 4** (note the exaggerated scale in the Y-direction). Total time of flight is $t_f = 289$ ms and the particle impacts 6.2 km east of the field line intersection point with the Earth.

If the distance traveled by the particle during one cyclotron orbit L_c is small in comparison to the gradient length scale of the magnetic field L_B , then we can assume that the magnetic moment μ is conserved to all orders. This ratio is largest during particle injection, when the ambient field strength is weakest. From Equation (2b), this ratio is computed as,

$$\epsilon = \frac{L_c}{L_B} = \frac{v_0}{\omega_c} \left/ \frac{B}{dB/dR} \right| = \frac{3m_0\gamma\beta c R_0^2}{qD} \quad (6)$$

where m_0 is the rest mass of the electron and q is the fundamental charge.

For these injection conditions, we have $\epsilon \approx 6.7 \times 10^{-3}$ and can therefore safely assume conservation of μ . Departure of the particle trajectory from the magnetic field line in the $-\hat{\mathbf{y}}$ direction can therefore be accurately described by single-particle-motion drifts. The particle initially drifts away from the field line and then appears to return due to the radial convergence of the

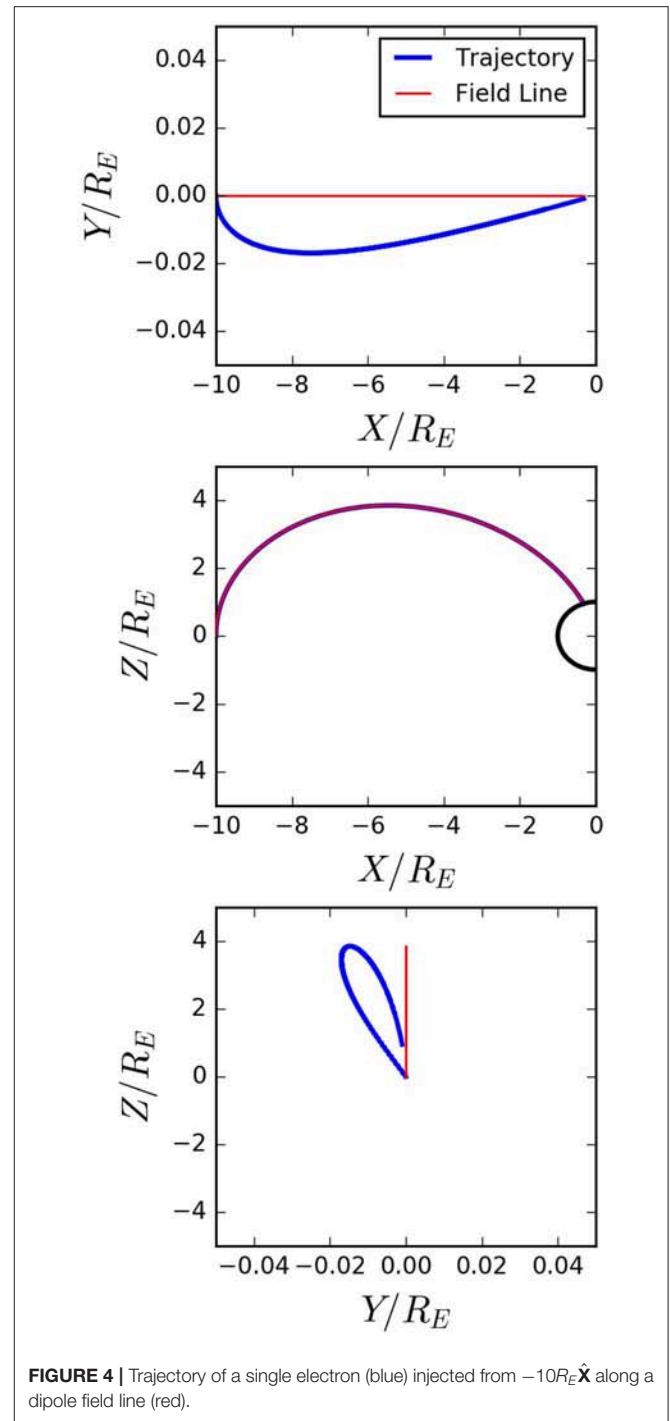


FIGURE 4 | Trajectory of a single electron (blue) injected from $-10R_E\hat{\mathbf{x}}$ along a dipole field line (red).

field lines approaching the Earth. For a relativistic electron, the drift velocity due to curvature and ∇B drifts is given by,

$$\mathbf{v}_d = \frac{\gamma m_0}{2q} \left(2v_{\parallel}^2 + v_{\perp}^2 \right) \frac{\mathbf{B} \times \nabla B}{B^3} \quad (7)$$

where $v_{\parallel} = v_0 \cos \delta$ and $v_{\perp} = v_0 \sin \delta$ are the parallel and perpendicular velocity components, respectively, and with reference to the local magnetic field.

Assuming that $v_{\perp} \ll v_{\parallel} \approx v_0$, using the field line geometry from Equation (2b) and integrating over the field line path, the total displacement of the particle impact location from the field line intersection point can be approximated as,

$$\Delta D \approx 1.131 \frac{m_0 \gamma \beta c R_E^3}{qD} \left(\frac{R_0}{R_E} \right)^{3/2} \quad (8)$$

For our reference beam conditions $\Delta D \approx 5.9 \text{ km}$, which is in reasonable agreement with the simulation.

3.2. Change in Impact Location With Magnetospheric Conditions

To explore the influence of magnetospheric conditions on the impact location of the beam, we simulate the injection of an electron into realistic BATS-R-US magnetospheric field geometry. The test case is the St. Patrick's day magnetospheric storm over the 17th and 18th of March 2015 (Jacobsen and Andalsvik, 2016). Simulations were run for seven different magnetospheric conditions, encompassing the pre-storm, Interplanetary Shock (IPS) arrival, storm main and recovery phases. Particles are injected from the equatorial plane at $-5R_E \hat{\mathbf{x}}$ on the midnight side of the noon-midnight meridian. Injection from this distance results in a favorable probability of injection into the loss-cone (see Willard et al., 2019 for further details).

Figure 5 shows the trajectory of each electron, which closely match that of their respective field lines. The final impact locations of each pulse are converted from geomagnetic coordinates to true longitude and latitude coordinates via the International Geomagnetic Reference Field (Thébault et al., 2015). The impact longitude is then shifted by some reference value to indicate impact over the North American continent, **Figure 6** shows the impact location of each of the beams. The separation of the impact positions clearly demonstrates that different phases of the storm will be highly distinguishable due to several-degree latitude separation between impact locations. Results from section 4 will show that these shifts are many orders of magnitude larger than the beam spot size at the top of the atmosphere and therefore will be clearly separable in ground station measurements. It is important to point out that ground stations will clearly need to be placed to cover observation over very large areas, particularly during intense geomagnetic activity when impact locations can be separated by more than 2,000 km.

In this idealized hypothetical diagnostic campaign, it is assumed that the satellite will be capable of injecting particles from an identical location at all times. Although this scenario provides clarity regarding Question 2, as posed in the introduction, it will most likely not be the case in reality. While a more thorough investigation of orbits is ongoing, possible realistic orbits could include geosynchronous, or sun-synchronous orbits, with perigee and apogee ranging between $5 - 10R_E$, allowing for multiple injection radii to be sampled.

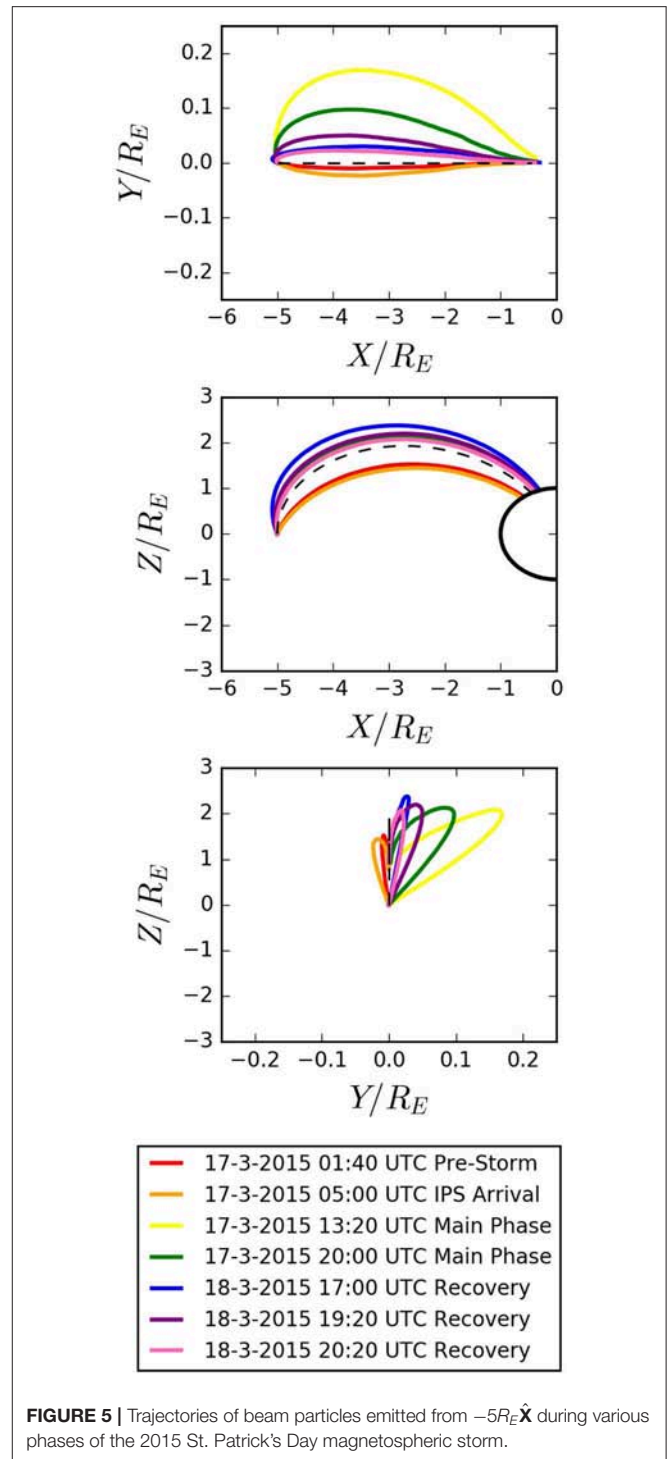


FIGURE 5 | Trajectories of beam particles emitted from $-5R_E \hat{\mathbf{x}}$ during various phases of the 2015 St. Patrick's Day magnetospheric storm.

4. EVOLUTION OF BEAM CROSS-SECTION DENSITY PROFILE

In this section we consider the third question posed in the introduction to this paper, whether the beam will remain sufficiently focused during propagation. This can be determined by studying the evolution of the beam cross-section density

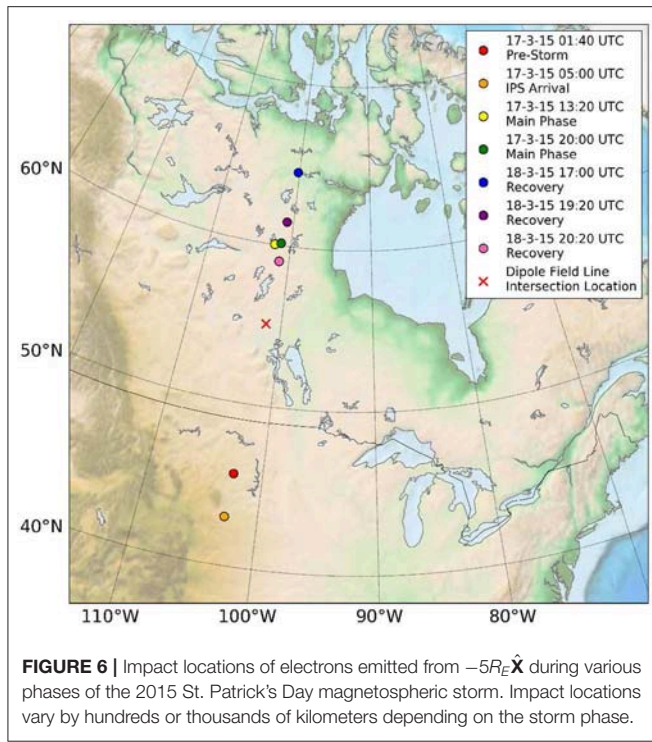


FIGURE 6 | Impact locations of electrons emitted from $-5R_E \hat{x}$ during various phases of the 2015 St. Patrick's Day magnetospheric storm. Impact locations vary by hundreds or thousands of kilometers depending on the storm phase.

profile from injection to impact with the Earth. The evolution of the beam density profile will be affected by beam initial conditions; energy, energy spread, emittance, beam radius and injection angles, as well as the geometry of the magnetospheric field lines.

Significant headway can be made in this analysis by considering a simple ensemble of electrons moving in a uniform magnetic field. Let the electrons evolve in the $\hat{x} \times \hat{y}$ plane with magnetic field in the $\hat{z}$ direction. Each particle has initial velocity vector $\mathbf{v}_\perp(t=0) = (v_\perp + \delta v_x)\hat{x} + \delta v_y\hat{y} + v_\parallel\hat{z}$ with δv_x and δv_y randomly sampled from a two-dimensional Maxwellian distribution with RMS velocity $\langle v_\perp \rangle$ given by the beam emittance ε_r . From Newton's second law and the Lorentz force law, particle positions will evolve as,

$$\begin{aligned} x(t) &= \frac{v_\perp}{\omega_c} \sin(\omega_c t) + \frac{\delta v_x}{\omega_c} \sin(\omega_c t) - \frac{\delta v_y}{\omega_c} \cos(\omega_c t) \\ y(t) &= \frac{v_\perp}{\omega_c} \cos(\omega_c t) + \frac{\delta v_x}{\omega_c} \cos(\omega_c t) + \frac{\delta v_y}{\omega_c} \sin(\omega_c t) \end{aligned} \quad (9)$$

where $\omega_c = qB/\gamma m_0$ is the relativistic electron cyclotron frequency.

The first component of each equation describes simple cyclotron motion with radius $r_c = v_\perp/\omega_c$. The remaining two components of each equation describe the evolution of the beam RMS radius r_b , also known as the beam envelope,

$$r_b = \sqrt{\langle \delta v_x \rangle^2 + \langle \delta v_y \rangle^2} = \frac{\langle v_\perp \rangle}{\omega_c} (1 - \cos(\omega_c t)) \quad (10)$$

The envelope of the ensemble of particles is therefore expanding and contracting on the cyclotron time scale and in phase with the centroid motion of the entire beam rotating at radius r_c . In sections 4.1 and 4.4 below, we show that for conditions near the reference values we generally have $r_c > r_b$, and therefore, shortly after injection the beam density profile evolves similar to that shown in **Figure 7A**.

Since there is a spread in beam energy the cyclotron frequencies of each particle will differ slightly due to the Lorentz factor γ in the denominator. The RMS spread in the cyclotron frequency ω_c is approximately related to the energy spread via,

$$\Delta\omega_c \approx \frac{q^2 B E_0}{(\gamma m c)^2} \Delta E \quad (11)$$

In section 4.3, we show that for near beam reference conditions, and injection from $10R_E$, the particles will spread many periods in gyro-phase during their time of flight and that the beam density profile will evolve into that of a corkscrew as their gyro-phase and position along the beam are spread. At impact with the atmosphere, this corkscrew will be projected into a ring, as shown in **Figure 7B**.

In the following sections, the relative magnitudes of the centroid cyclotron radius r_c and beam envelope radius r_b are computed via a more complete analysis. This analysis includes the effects of finite mini-pulse size, finite pulse length, emittance, energy spread, beam self-forces and magnetic field geometry. Without loss of generality, we can continue to decouple the evolution of the beam centroid motion from that of the beam envelope (Qin et al., 2010, 2011; Chung and Qin, 2018). We thus proceed by considering the evolution of the beam centroid and the gyro-phase spread due to energy spread. We then study the evolution of the beam envelope and show that the final envelope size is generally smaller than the beam cyclotron radius. Next, we consider the optimum injection angle for a beam and derive restrictions on the pointing accuracy of the satellite, as well as the loss fraction of particles for beams with large emittance. Finally, we incorporate the motion of the satellite into our calculations and compare these results to ballistic simulations.

4.1. Evolution of Beam Mini-Pulse Centroid

Evolution of the beam centroid is modeled by the dynamics of a single electron injected into an ideal dipole magnetosphere. It is assumed that the electron has properties near to those of the beam reference conditions. Since the particle is relativistic, higher order terms of the asymptotic expansion must be considered when computing the magnetic moment (Porazik et al., 2014). For an axisymmetric field and a particle injected from the geomagnetic mid-plane, a second order expression for the magnetic moment at injection is given by (Gardner, 1966),

$$\mu \approx \mu^{(0)} + \mu^{(1)} + \mu^{(2)} \quad (12a)$$

$$\mu^{(0)} = \frac{m v_\perp^2}{2B} \quad (12b)$$

$$\mu^{(1)} = -\frac{m^2 B'}{2qB^3} (v_0^2 + v_\parallel^2) v_\lambda \quad (12c)$$

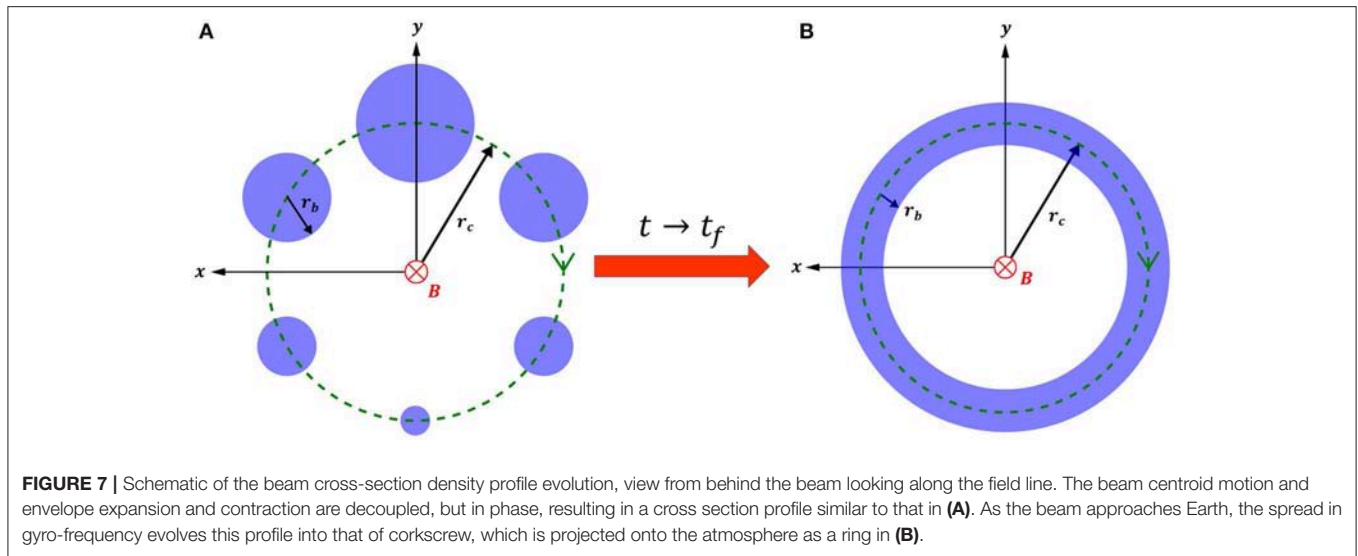


FIGURE 7 | Schematic of the beam cross-section density profile evolution, view from behind the beam looking along the field line. The beam centroid motion and envelope expansion and contraction are decoupled, but in phase, resulting in a cross section profile similar to that in (A). As the beam approaches Earth, the spread in gyro-frequency evolves this profile into that of corkscrew, which is projected onto the atmosphere as a ring in (B).

$$\mu^{(2)} = \frac{m^3}{2q^2} \left\{ \frac{B'^2}{B^5} \left[\frac{1}{2} (3v_\lambda^2 + v_\parallel^2) (v_0^2 + v_\parallel^2) + \frac{3}{8} v_\perp^4 \right] - \frac{B''}{2B^4} \left[v_\lambda^2 v_\parallel^2 + \left(v_\lambda^2 + \frac{1}{4} v_\perp^2 \right) (v_0^2 + v_\parallel^2) \right] + \frac{B'}{2rB^4} \left[v_\lambda^2 v^2 - v_\perp^2 v_\parallel^2 - \frac{5}{4} v_\perp^2 (v_0^2 + v_\parallel^2) + 2v_\lambda^2 v_\parallel^2 \right] \right\} \quad (12d)$$

Where $v_\perp = v_0 \sin \delta$, $v_\lambda = v_\perp \sin \lambda$, $v_\parallel = v_0 \cos \delta$ and $B' = dB/dR$, and angles δ and λ are illustrated in **Figure 2**. In a dipole field, and when $\delta \ll 1$, such that $v_\lambda \leq v_\perp \ll v_\parallel \approx v_0$, Equation (12) can be approximated as,

$$\mu \approx \frac{m_0 \gamma v_\perp^2 R_0^3}{2D} \left(1 - 2\epsilon \frac{v_0}{v_\perp} \sin \lambda + \epsilon^2 \frac{v_0^2}{v_\perp^2} \right) \quad (13)$$

Where ϵ is defined in Equation (6). Note that $\mu \neq 0$ when the particle is injected directly along the field line. With non-zero magnetic moment, the particle will experience magnetic mirroring forces, and therefore a transition from parallel to perpendicular kinetic energy as it moves into an increasing magnetic field strength closer to the Earth. Consider again the example simulation of section 3.1, for a particle injected directly along the field line ($\delta = 0^\circ$) from $-10R_E \hat{\mathbf{x}}$ and with reference beam properties, at impact 5.7% of the initially parallel kinetic energy is converted into perpendicular kinetic energy, resulting in a cyclotron radius of $r_c = 21.8 \text{ m}$.

Although not discussed in detail here, it should be noted that for a particle injected from this location with sufficiently high energy $E_0 \gtrsim 4 \text{ MeV}$, the particle loss cone may no longer include injection directly along the field line (Porazik et al., 2014).

At impact with the Earth, the magnetic field is strong enough and the perpendicular velocity large enough, that the most significant contribution to the magnetic moment is given by the zeroth order component. Relating the initial magnetic moment, to the final moment at the Earth, $\mu_E^{(0)} = \mu$ gives a general relationship for the final cyclotron radius at impact for any particle injected from the geomagnetic equatorial plane onto a dipole field line,

$$r_c = \sqrt{\frac{m_0^2 \gamma^2 \beta^2 c^2 R_E^3}{q^2 D B_E} (\sin^2 \delta - 2\epsilon \sin \delta \sin \lambda + \epsilon^2)} \quad (14)$$

where $B_E = B(R_E, \theta_E)$ is the magnetic field strength at the field line intersection point with the Earth.

For $\delta = 0^\circ$, injection from $-10R_E \hat{\mathbf{x}}$ and reference beam properties, Equation (14) predicts a final cyclotron radius of 21.7 m in near identical agreement to simulations. **Figure 8** shows how the final cyclotron radius of this simulation is predicted to vary with injection radius, injection energy and injection angles δ (with fixed $\lambda = -90^\circ$), and λ (with fixed $\delta = 0.5^\circ$). These results are verified via ballistic simulations.

4.2. Optimum Injection Angle and Pointing Accuracy

Figures 8C,D suggest that there is an optimum injection angle which will minimize the magnetic moment, and therefore final beam centroid cyclotron radius. **Figure 8D** and Porazik et al. (2014) indicate that the optimum azimuthal injection angle is $\lambda = -\pi/2$. The optimum polar injection angle can be obtained by setting Equation (14) equal to zero, which for small angles gives solution $\delta \approx \epsilon$. Therefore $(\lambda, \delta) = (-\pi/2, \epsilon)$ describe a pair of injection angles which yield zero cyclotron radius at impact with the upper atmosphere.

The limits of the loss cone in a dipole field with $\lambda = -\pi/2$ can be described by Porazik et al. (2014),

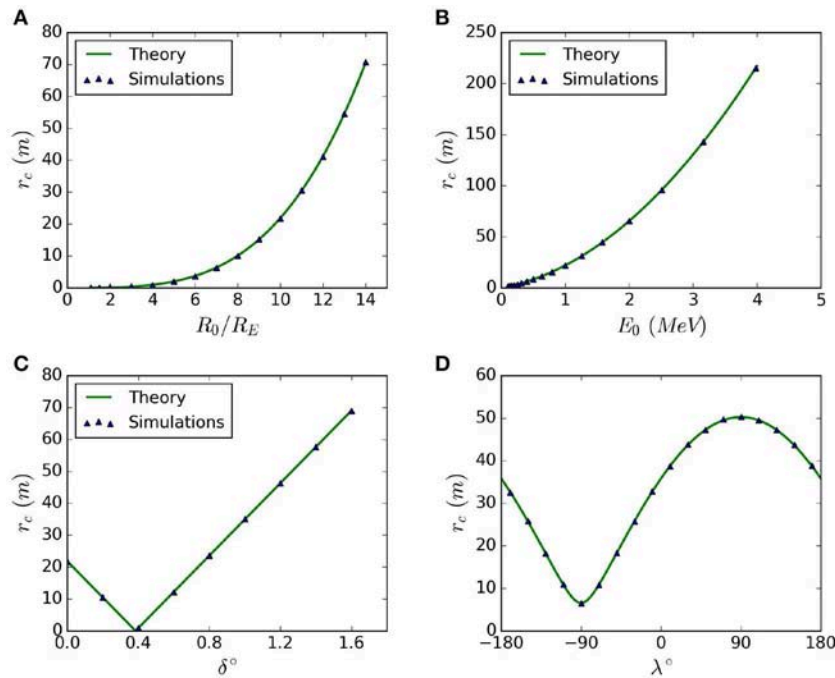


FIGURE 8 | Variation of beam cyclotron radius r_c at impact with the Earth for the reference case ($\delta = 0^\circ$, injection from $-10R_E\hat{\mathbf{x}}$ and reference beam properties of **Table 1**) with changes in **(A)** initial injection radius R_0 , **(B)** initial injection energy E_0 , **(C)** with polar injection angle δ for fixed $\lambda = -90^\circ$, **(D)** with azimuthal injection angle λ for fixed $\delta = 0.5^\circ$.

$$\sin^2 \delta - \frac{\epsilon}{4} (5 \sin \delta + \sin 3\delta) + \frac{\epsilon^2}{384} (275 + 68 \cos 2\delta + 41 \cos 4\delta + 4(43 + \cos 2\delta) \sin^2 \delta) = \frac{(R_E/R_0)^3}{\sqrt{4-3(R_E/R_0)}} \quad (15)$$

Using the small angle approximation, and making the substitution $\delta \rightarrow \delta_i - \epsilon$ we obtain,

$$\delta_i^2 \approx \frac{(R_E/R_0)^3}{\sqrt{4-3(R_E/R_0)}} \quad (16)$$

Where δ_i is the injection angle with respect to the frame shifted by angles $(-\pi/2, \epsilon)$ from the local magnetic field vector. In this frame we recover the traditional circular loss cone, with angular radius defined by δ_i from Equation (16). We can consider δ_i as a minimum bound for the pointing accuracy of the satellite mounted electron accelerator. For standard beam conditions $\epsilon = 0.38^\circ$ and $\delta_i \approx 1.31^\circ$.

The limits on injection angle may also set restrictions for the radial beam emittance ϵ_r , which determines the initial radial velocity spread at beam injection. For a sufficiently large ϵ_r , a significant portion of the injected particles may be injected outside of the loss cone, reducing the signal observed at the top of the atmosphere. For the standard beam conditions, however, we have an RMS spread in injection angles described by $\Delta\delta \approx \epsilon_r/r_{b,i} = 0.03^\circ$, therefore if fired along the optimum injection angle beam particles will remain well inside the loss cone.

4.3. Spreading of Particle Gyro-Phase Due to Energy Spread

As the particles stream along field lines their gyro-orbits will decorrelate due to a spread in their Lorentz factors γ , and therefore gyro-frequencies. The RMS shift in gyro-phase ψ of a particle with respect to initial energy E_0 can be computed from,

$$\frac{d\psi}{dS} = \frac{\Delta\omega_c}{\beta c} = \frac{q^2 E_0 \Delta E}{\beta c^3 \gamma^2 m_0^2} B(S) \quad (17)$$

Where S is the arc-length of the beam measured from injection and $\Delta\omega_c$ is given by Equation (11).

Equation (17) is integrated from injection until impact with the Earth via the ODE integration package LSODE (Hindmarsh, 1980), implemented in Python with SciPy (Oliphant, 2007). The field strength $B(S) = B(r(S), \theta(S))$ is adjusted at each integration point via Equation (2b). For reference beam conditions, **Figure 9** shows how the RMS gyro-phase spread changes with time of flight. The rate of phase shift increases closer to the Earth as the particle transits a steeper magnetic field gradient, resulting in a total phase spread of ≈ 50 gyro-periods. Therefore, the initial centroid motion of the beam mini-pulse will transition into a rotating corkscrew, which projects into a ring at impact with the Earth (see **Figure 7**).

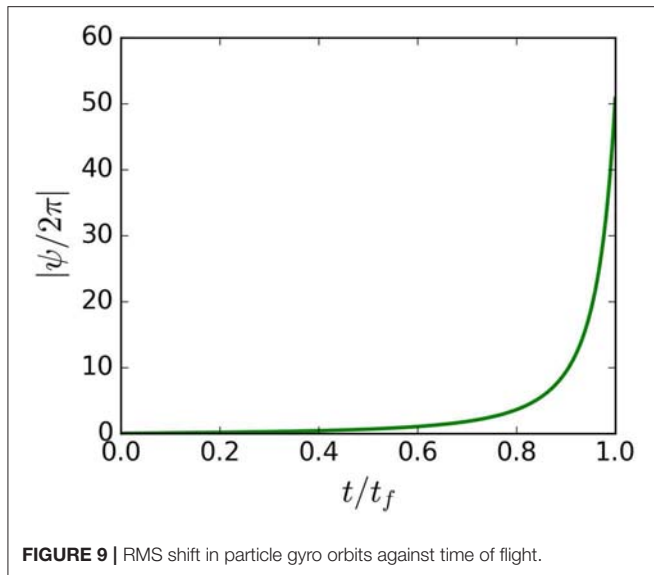


FIGURE 9 | RMS shift in particle gyro orbits against time of flight.

4.4. Evolution of Beam Mini-Pulse Envelope

The beam envelope evolves due to beam initial conditions, self-generated electromagnetic forces, and the applied magnetic field strength. At maximum expansion, the beam mini-pulse length remains $\sim 10^3$ times longer than the radius, therefore a mini-pulse is modeled as an infinitely long beam. The mini-pulse envelope, r_b evolves according to the one-dimensional beam envelope equation (Reiser, 2008),

$$\frac{d^2 r_b}{ds^2} = -k_0^2 r_b + \frac{K}{r_b} + \frac{\varepsilon_r^2}{r_b^3} \quad (18)$$

Where $k_0 = \frac{qB(S)}{2m_0 c \gamma \beta}$ is due to focusing from the applied magnetic field and $K = \frac{qI_0}{2\pi \varepsilon_0 m_0 c^3 \beta^3 \gamma^3}$ is the perveance, which captures the influence of beam self-charge and self-magnetic field, and ε_r is the radial emittance.

Equation (18) is integrated via the same techniques described in section 4.3. Since the ballistic propagation code does not incorporate the effects of space charge, Equation (18) is also solved for the case without perveance ($K = 0$) to allow for comparison with ballistic simulations.

The solution to the ODE predicts an oscillating beam envelope, however due to gyro-phase mixing the particles will likely fill out a profile at the extrema of these oscillations. The solution to the beam envelope equations with no perveance, and the extrema profiles for the case with and without perveance are shown in **Figure 10**. Initially the beam radius blows out to a size on the order of hundreds of meters, and then as it propagates toward the Earth, the increasing magnetic field strength focuses the beam. For the reference conditions, the final beam radius at impact with the Earth is 2.6 m with perveance and 2.2 m without perveance.

Equation (18) is solved parametrically to determine the influence on final envelope radius $r_{b,f}$ for changes in initial energy

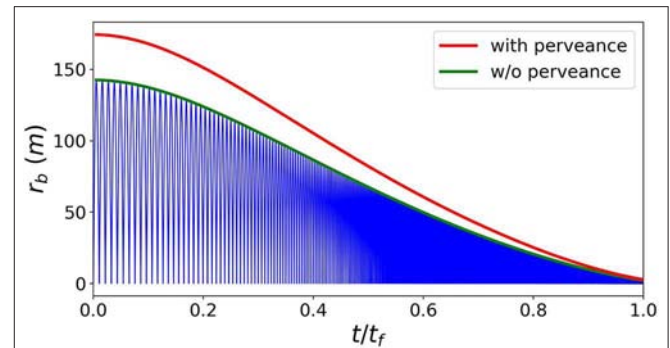


FIGURE 10 | Solution to the envelope equations without beam perveance (blue lines). Extrema solutions with (red) and without perveance (green) for the reference beam conditions.

E_0 , injection radius R_0 , radial emittance ε_r , initial beam radius $r_{b,i}$ and beam current I_0 . The results of these parameter scans for cases with and without perveance are shown in **Figure 11** along with the corresponding final cyclotron radius r_c for these conditions.

For the reference beam current of 1 mA, the final beam envelope radius is only weakly influenced by beam perveance. This is due to the small magnitude of the average current as well as the self-generated magnetic field which acts to cancel out a large fraction of the beam self-charge for $\gamma \approx 3$. This demonstrates that despite self-forces being neglected, the use of ballistic simulations is suitable for modeling beams with similar properties to those here.

Other observable trends include that increasing the initial beam energy E_0 results in a larger final radii, since the increased electron momentum increases the particle cyclotron radius for the same applied magnetic field. Increasing the beam injection radius R_0 similarly results in an increased final beam radius, due to the weaker magnetic field and therefore larger initial cyclotron radius at injection. Unsurprisingly, increasing beam radial emittance ε_r increases final beam radius since the particles are initialized with a larger RMS perpendicular velocity. Increasing the initial beam radius $r_{b,i}$ results in a smaller final beam radius since the initial current density (and therefore self-electric fields) are reduced. Finally, increasing beam current I_0 results in an increased final radius due to the increased current density at beam initialization. Therefore, the larger the beam current, the less suitable ballistic simulations become for modeling the beam.

It is clear that for conditions near the reference beam properties, the cyclotron radius of the beam centroid dominates the profile of the final particle density distribution since $r_{b,f}/r_c \ll 1$. The cyclotron radius r_c is therefore the most important quantity when considering final beam spot size. **Figure 11** shows that for increasing beam energy E_0 and injection radius R_0 , this ratio will become even smaller. Only for large increases to ε_r , I_0 or decreases to $r_{b,i}$ will this ratio be larger than unity, and then the evolution of the beam envelope may become a more important consideration.

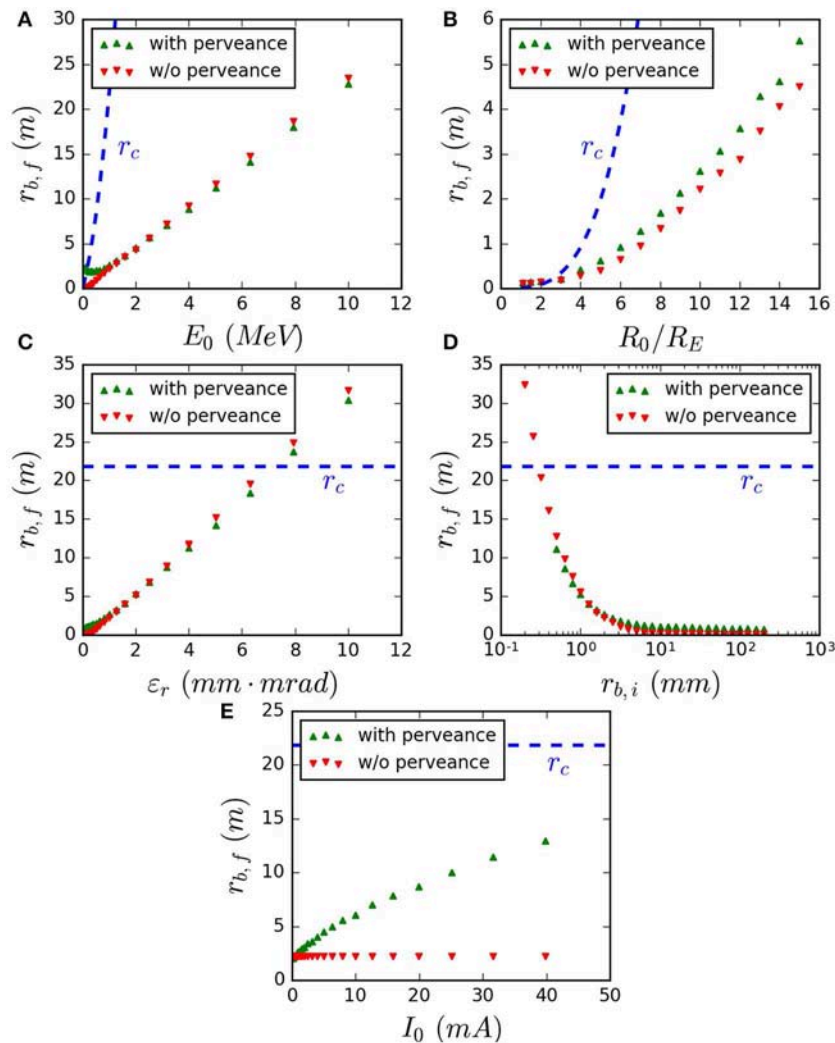


FIGURE 11 | Final RMS beam envelope radius at impact with the Earth with parameter scans in **(A)** beam energy E_0 , **(B)** injection radii R_0 , **(C)** radial emittance ϵ_r , **(D)** initial beam radius $r_{b,i}$, **(E)** injection current I_0 . Blue dashed lines show the equivalent cyclotron radius for each condition as per Equation (14).

4.5. Effects of Satellite Motion on Beam Pulse Impact Distribution

A single electron beam pulse (see **Figure 3**) consists of 100 mini-pulses and total pulse time of $T_p = 0.5$ s. At these time scales it becomes important to consider the motion of the electron gun platform. If this platform were to remain stationary during firing, then the impact density distribution would appear similar to that of a single mini-pulse; however, since the accelerator is attached to a moving satellite, the beam impact location will be smeared out, as each mini-pulse is injected onto a slightly different field line.

Assuming that the satellite is in a circular equatorial orbit, a simple approximation for the satellite angular velocity Ω_0 can be obtained using Newton's second law and equating the gravitational force of the Earth and the centripetal force exerted by the orbiting satellite,

$$\Omega_0 = \frac{V_0}{R_0} = \sqrt{\frac{GM_E}{R_0^3}}, \quad (19)$$

where G is Newton's gravitational constant and M_E is the mass of the Earth. An equatorial orbit will likely be prograde therefore when calculating the beam impact spread on the surface of the Earth we can subtract the angular velocity of the Earth itself; $\Omega_E = 2\pi/T_E$, where $T_E = 86,400$ s is the period of the Earth's rotation. Therefore the azimuthal shift of the pulse impact location is simply $\Delta\phi = (\Omega_0 - \Omega_E)T_p$. Since the dipole field is cylindrically symmetric, the beam spread must be calculated using the cylindrical radius $R_E \sin \theta_E$ at the impact location of the dipole field line with the Earth. Therefore, the total shift in impact location Δd is given by,

$$\Delta d = R_E \Delta\phi \sin \theta_E = T_p R_E \left(\frac{\sqrt{GM_E R_E}}{R_0^2} - \frac{2\pi}{T_E} \sqrt{\frac{R_E}{R_0}} \right) \quad (20)$$

Figure 12 shows the pulse impact location displacement against pulse time for various initial injection radii R_0 , where **Figure 12A** shows the case without the Earth's rotation and **Figure 12C**

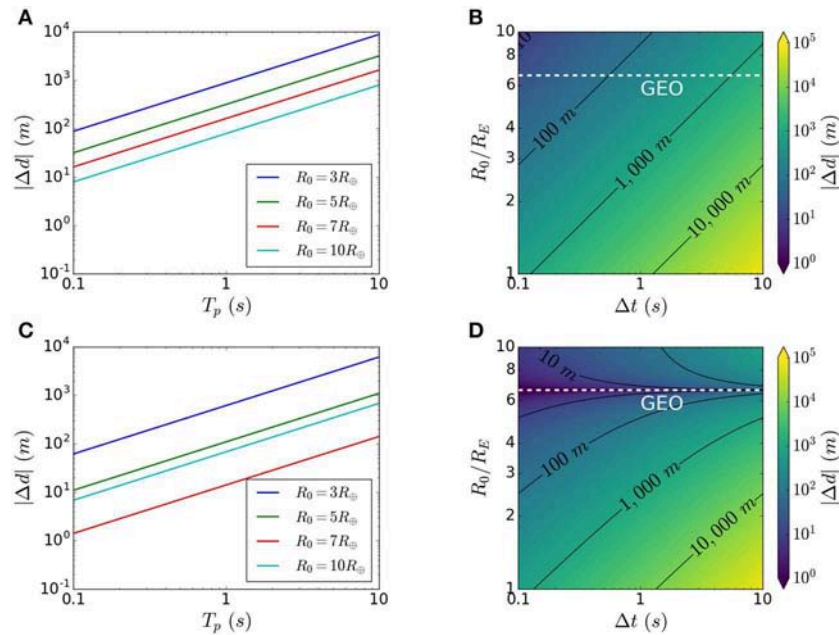


FIGURE 12 | Displacement of the center of the beam spot $|\Delta d|$ due to motion of a satellite-borne particle accelerator **(A,B)** satellite motion only, **(C,D)** satellite motion and Earth rotation. The white dashed line indicates the radius corresponding to Geostationary Orbit (GEO).

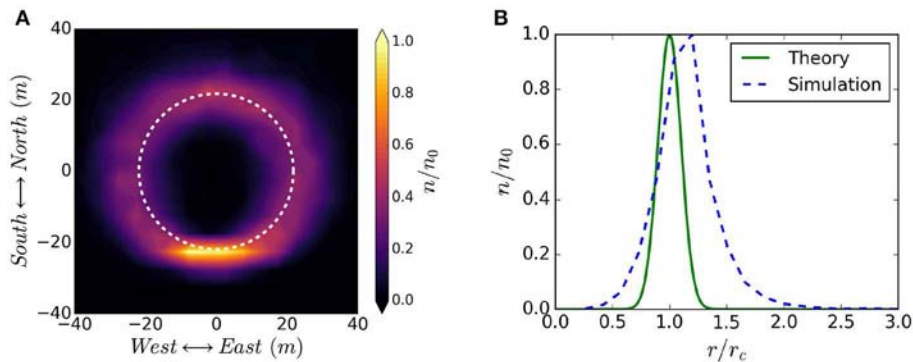


FIGURE 13 | (A) Normalized electron density distribution at impact with the Earth, the theoretically predicted cyclotron radius from Equation (14) is overlaid as a dashed white line. **(B)** Normalized radial electron density distribution compared with theoretical predictions for beam cyclotron radius (distribution mean) and final beam envelope RMS radius (distribution spread). An animation of beam particle evolution from injection until impact with the atmosphere for these results can be viewed here: https://youtu.be/ZupUFIF2_yE.

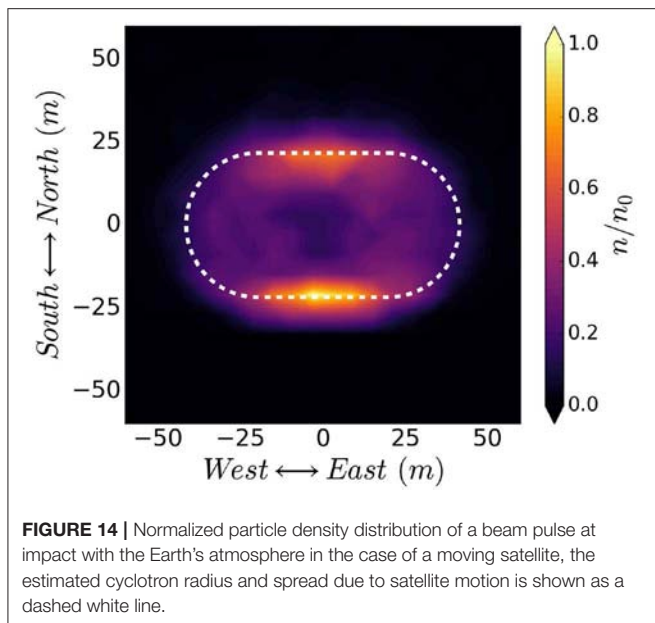
includes the Earth's rotation. **Figures 12B,D** show heat maps of the same situation for injection radii under consideration. For the beam reference conditions and injection from $10R_E$, the spread of the center of the beam is 79 m .

4.6. Comparison With Ballistic Simulations

To explore the validity of the above results we simulate a single mini-pulse of electrons injected from $10R_E$ along a dipole field line. An animation of 200 particles sampled from this simulation can be found here https://youtu.be/ZupUFIF2_yE, and **Figure 13A** shows the normalized density distribution of particles impacting the upper atmosphere (in the $\hat{\theta} \times \hat{\phi}$

plane). As expected, the impact distribution is ring-shaped rather than circular. The dashed white line shows the predicted final cyclotron radius $r_c = 21.8\text{ m}$ from Equation (14). Despite decorrelation of particle gyro-orbit phase, the brighter region at the bottom of the ring demonstrates that a large number of the particles remain closely correlated. In reality, the beam perveance will modify the oscillation frequency of the beam envelope, and therefore if beam self-forces were included we would expect a more uniform distribution of density around the ring.

Figure 13B shows the radial electron density distribution (measured with respect to the origin in **Figure 13A**) and the predicted density distribution given by $n/n_0 =$



$\exp((r - r_c)^2/2r_{bf}^2)$. The mean location of these two profiles show agreement within 15% between the final cyclotron radius of the simulation and the prediction of Equation (14). There is less clear agreement between the simulated and predicted RMS beam envelope radii, with the simulation showing an envelope radius approximately double that of the 2.2 m predicted by solving Equation (18). A possible explanation for this discrepancy is that the energy spread ΔE results in additional beam broadening in the East to West direction due to the energy dependence of the ∇B and curvature drifts. Particles with higher energy will drift further to the East than those with lower energy, resulting in a smearing of the density profile. This may also explain the small discrepancy between the predicted and observed average cyclotron radii.

For completeness we have also produced an animation demonstrating reflection of the beam due to magnetic mirroring. This simulation is for identical conditions, except with injection angles $\delta = 2^\circ$, $\lambda = -90^\circ$ and can be found here <https://youtu.be/Y-amwRDZruo>.

Next, 100 mini-pulses, which comprise a single electron pulse, are injected with beam reference properties along a dipole field line from $10R_E$. The final normalized impact density distribution is shown in **Figure 14**. The white dashed line shows the theoretically predicted cyclotron radius combined with the drift predicted due to satellite motion from Equation (20) (without Earth rotation). The center of the spot appears filled out when compared to that of the mini-pulse in **Figure 13**, and is caused by smearing of the impact ring due to satellite motion. The effect of the true satellite motion and relative rotation of the Earth can be incorporated for specific orbital parameters if required. Details of beam mini-pulse overlapping can be found in the **Appendix** to this publication.

At impact with the atmosphere, relativistic electrons produce very-low frequency (VLF) waves due to secondary ionization, optical emissions due to excitation of neutrals and high energy

photons due to bremsstrahlung (Marshall et al., 2014). If their signature is strong enough, both VLF waves and optical spectra can be detected by ground stations, and high energy photons may be observed by high altitude or orbital observatories.

Beam properties at the top of the atmosphere can be used as initial conditions for Monte-Carlo collision models, such as those developed in Marshall et al. (2014), to determine the emission profile of the beam interacting with the atmosphere. Optical emission occurs mostly within the D region of the atmosphere once the beam has spread out to hundreds of meters in radius. The optical photon flux is therefore relatively insensitive to beam radius at the top of the atmosphere, provided $r_{bf} \lesssim 100$ m, which is satisfied for the reference beam conditions. Specific ground station sensors can then be considered to determine a signal-to-noise ratio, and therefore whether the beam will be detectable.

Marshall et al. (2014) shows that for similar beam energies and fluxes referred to in this paper, the resulting emission spectra will produce significant and detectable signatures. Private communications between users of these tools confirm that this is the case, and further details can be found in Sanchez et al. (in preparation).

5. CONCLUSION

This paper explores the dynamics of an electron beam propagating from injection into the magnetosphere until impact with the Earth. Injected from the geomagnetic equatorial plane along a dipole field line, particles were found to shift from their original field line due to single-particle-motion drifts. The total integrated drift motion is on the order of kilometers, and therefore when compared to the radius of the Earth the particle impact location is nearly identical to that of the field line intersection point.

Particles were injected from an identical location at 5 Earth radii during different phases of a simulated magnetospheric storm. The phase of the storm was found to strongly influence the impact location of the particles, shifting them hundreds to thousands of kilometers. This simulated diagnostic campaign demonstrates that even for injection radii near to geostationary orbit, magnetospheric weather can have a large observable influence on the impact location of a beam propagating from a satellite down into the atmosphere. It also demonstrates that a wide ground station coverage area will be required to detect these signatures.

Evolution of the beam cross-section was studied by considering the separate dynamics of the beam centroid motion and evolution of the beam envelope. For beam properties near those considered in this work, the final beam centroid cyclotron radius was found to be the most important parameter when estimating the beam spot size at the top of the atmosphere. For the provided beam reference conditions, a single beam mini-pulse impacts with a density profile in the shape of a ring, with major radius ~ 22 m, and ring thickness ~ 2 m. When considering a single pulse (multiple mini-pulses) the beam spot size is additionally spread 10s to 100s of meters due to motion of the orbiting

electron accelerator. With sufficient pointing accuracy and with reference beam emittance, it is possible to inject beam particles into the loss cone. It is shown that the beam spot size will remain sufficiently narrow to allow detection from ground stations on the surface of the earth. Future work will explore how beam-plasma instabilities may modify the final beam spot size.

Demonstrating theoretically that the beam will be detectable by ground stations, and that magnetospheric events will significantly influence the beam impact location provides validation for two of the key requirements of this proposed diagnostic.

DATA AVAILABILITY STATEMENT

The datasets generated for this study are available on request to the corresponding author.

AUTHOR CONTRIBUTIONS

AP contributed significantly to original research, and prepared the manuscript. PP contributed significantly to original research in this article. MG-M, KA, and DS contributed to original research for this article. IK, JJ, and ES are principal investigators for this research, they contributed and guided this research.

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SUPPLEMENTARY MATERIAL

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Relativistic Particle Beams as a Resource to Solve Outstanding Problems in Space Physics

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The Sun's connection with the Earth's magnetic field and atmosphere is carried out through the exchange of electromagnetic and mass flux and is regulated by a complex interconnection of processes. During space weather events, solar flares, or fast streams of solar atmosphere strongly disturb the Earth's environment. Often the electric currents that connect the different parts of the Sun-Earth system become unstable and explosively release the stored electromagnetic energy in one of the more dramatic expressions of space weather—the geomagnetic storm and substorm. Some aspects of the magnetosphere-ionosphere connection that generates auroral arcs during space weather events are well-known. However, several fundamental problems remain unsolved because of the lack of unambiguous identification of the magnetic field connection between the magnetosphere and the ionosphere. The correct mapping between different regions of the magnetosphere and their foot-points in the ionosphere, coupled with appropriate distributed measurements of plasma and fields in focused regions of the magnetosphere, is necessary to establish unambiguously that a given magnetospheric process is the generator of an observed arc. We present a new paradigm that should enable the resolution of the mapping ambiguities. The paradigm calls for the application of energetic electron beams as magnetic field tracers. The three most important problems for which the correct magnetic field mapping would provide closure to are the substorm growth phase arcs, the expansion phase onset arcs and the system of arcs that emerge from the magnetosphere-ionosphere connection during the development of the early substorm expansion phase phenomenon known as substorm current wedge (SCW). In this communication we describe how beam tracers, in combination with distributed measurements in the magnetosphere, can be used to disentangle the mechanisms that generate these critical substorm phenomena. Since the application of beams as tracers require demonstration that the beams can be injected into the loss cone, that the spacecraft potentials induced by the beam emission are manageable, and that sufficient electron flux reaches the atmosphere to be detectable by optical or radio means after the beam has propagated thousands of kilometers under

competing effects of beam spread and constriction as well as effects of beam-induced instabilities, in this communication we review how these challenges are currently being addressed and discuss the next steps toward the realization of active experiments in space using relativistic electron beams.

Keywords: relativistic beams, magnetic field mapping, magnetosphere-ionosphere coupling, storms and substorms, atmospheric effects of beams

INTRODUCTION

The Sun and the Earth are coupled in multiple ways. Heat coming from the Sun is the main energy source of Earth's weather. Solar ultraviolet emissions are the main source of the ionized layer of the Earth's upper atmosphere—the ionosphere. The particles and magnetic field emanated by the Sun flow over and merge with the Earth's magnetic field, reconfiguring it as well as the particle distributions trapped within. The Sun's interaction with the Earth's magnetic field, ionosphere and atmosphere leads to exchange of electromagnetic and mass flux which is regulated by a complex interconnection of processes (e.g., Vasyliunas, 1984; Lysak, 1990). Solar flares or fast streams of solar wind strongly disturb the Earth's surrounding environment known as the magnetosphere. Often during these events the electric currents that connect the different parts of the Earth's magnetosphere with the ionosphere become unstable and explosively release the stored electromagnetic and particle energy in one of the more dramatic expressions of space weather—geomagnetic storms and substorms (e.g., Gonzalez et al., 1994). These phenomena deposit large fluxes of energetic charged particles and electromagnetic energy into the atmosphere, driving the bright dynamic optical auroral displays. They also accelerate charged particles and inject them from regions deeper in the magnetotail into regions now thickly populated by commercial, scientific, and military spacecraft. The incident electromagnetic and particle fluxes can cause major ionospheric disturbances that impede communications and navigation during space weather events. The physical processes involved in substorms occur throughout the solar system and the universe: Substorms are observed on Saturn and Jupiter (Russell et al., 2000; Cowley et al., 2005; Mitchell et al., 2005; Kronberg et al., 2008) and the flares of energetic X-rays and gamma rays associated with such reconfigurations are observed routinely from our Sun and other stars (Masuda et al., 1994; Shibata, 1998).

Several fundamental questions about how the magnetosphere and the ionosphere are connected during storms and substorms remain unsolved. Understanding the connection is the most critical step toward understanding how the magnetosphere-ionosphere (M-I) system evolves from a stable to an unstable state. This will only be possible when we are able to unambiguously determine what processes and regions in the magnetosphere are linked to the aurora. This goal requires the ability to map the magnetic field lines that connect a given arc with its source region in the magnetosphere, and to measure the spatial and temporal evolution of the source region. The three most important questions for which the correct magnetic field mapping would provide closure to

are: How are the substorm growth phase arcs generated, how are the expansion phase onset arcs generated and how does the system of arcs and electric currents known as substorm current wedge (SCW; e.g., McPherron et al., 1973; Pytte et al., 1976) emerge during the early substorm expansion phase.

The Substorm Growth Phase

Ever since Akasofu and Chapman coined the term substorm (Akasofu and Chapman, 1961) and Akasofu described the auroral phenomenology of substorms (Akasofu, 1964), observational and theoretical investigations have been carried with the objective of explaining substorm evolution. These investigations have revealed that before the explosive release of energy occurs, there is an interval where energy from the Sun is being deposited in the Earth's magnetosphere. This interval is commonly referred to as the growth phase. The growth phase of the substorm is characterized by forming of multiple arcs, which brighten and remain stable for ~ 1 h (Figure 1; e.g., Akasofu, 1964; Lyons et al., 2002; Partamies et al., 2015). At the end of the growth phase, the M-I system reaches a state that allows explosive release of energy into the ionosphere, referred to as expansion phase. The start of the expansion phase is commonly referred to as breakup or expansion onset. Substorm auroral onset is characterized by a brightening near the equatorial boundary of the auroral oval, frequently along pre-existing growth phase arc (Akasofu, 1964). There are multiple unanswered questions regarding growth phase arcs and their relationship to the onset arc. For instance, it is unclear what makes the breakup arc different from all the others, aside from the obvious phenomenological evidence of brightening. It is therefore unclear how the electromagnetic coupling between the ionosphere and the magnetosphere differs for both sets of arcs. Recent modeling efforts of growth phase arcs carried out with the Rice Convection Model show the formation of a thin arc that extends several hours in magnetic local time in the transition region during the late growth phase, generated by large-scale adiabatic convection under equilibrium conditions (Yang et al., 2013). The arc in the pre-midnight sector is associated with precipitating electrons along an azimuthally elongated sheet of region-1 sense (whereby current flows into the ionosphere on the post-midnight side of the sheet and out of the ionosphere on its pre-midnight side) field-aligned current (FAC) just poleward of the main region-2 FAC, which has a polarity opposite that of the region-1 (Figure 2). The newly formed FACs are produced by a redistribution of pressure in the inner magnetosphere generated by convection and azimuthal particle drifts. As the

pressure redistribution continues, it concentrates Alfvén layers within progressively narrower L-shell range which maps to a narrower arc, thus forming what Yang et al. (2013) termed a “convection front.”

Models’ predictions for the growth phase electric current generators can be properly tested with an array of spacecraft in the nightside magnetosphere with the following attributes:

- 1) A main spacecraft emitting an electron beam to measure the unambiguous magnetic-field connection between the growth phase arcs observed with ground-based auroral cameras and the magnetospheric region where the beam was emitted from;
- 2) Three daughter spacecraft equidistant to the main spacecraft and to each other to enable measurement of plasma pressure and convection gradients in the radial and azimuthal directions. Variable radial separation of $\sim 10\text{--}4,000$ km among spacecraft will allow the *in-situ* testing of the growth phase models’ predicted pressure gradients in the magnetospheric equator.

The Substorm Expansion Phase

At the end of the growth phase, the M-I system reaches a state that allows explosive release of energy into the ionosphere, referred to as expansion phase. The transition usually occurs along an east-west-aligned auroral arc with a characteristic thickness between a few km (e.g., Hull et al., 2016) and ~ 30 km (e.g., Knudsen et al., 2001), emerges in the growth phase ~ 10 min before, and undergoes a sudden increase in brightness and subsequent rapid expansion azimuthally and longitudinally (e.g., Lyons et al., 2002; Hull et al., 2016, and references therein). Substorm current formation and evolution beyond growth phase is a process that involves both electron acceleration from static potentials at high altitude and Alfvénic acceleration mechanisms (e.g., Keiling, 2009 and references therein). Several case studies suggest that, as one of the most equatorward arcs intensifies during the transition from growth to onset of substorms, the arc may also develop filamentation into smaller scales which show wave properties (e.g., Wygant et al., 2002; Mende et al., 2003; Lessard et al., 2011; Hull et al., 2016). The transition is argued to be consistent with the notion that small-scale or dispersive Alfvén waves may be generated from larger-scale Alfvén waves and/or destabilization of macroscale currents (e.g., Chaston et al., 2011 and references therein).

As the magnetosphere-ionosphere system evolves into substorm breakup onset the brightening arc usually develops discrete rays, also called “beads,” pulsating in a wave-like form along the arc (Donovan et al., 2006; Liang et al., 2008; Henderson, 2009; Rae et al., 2010; Kalmoni et al., 2015; Nishimura et al., 2016; **Figure 3**). Onset waves have received attention because their optical properties seem to match at least some of the properties, such as growth rate and frequency, expected for several near-Earth instabilities that have been proposed as triggers of substorm expansion onset. These include cross-field current instabilities (Lui et al., 1991), various forms of fluid, hybrid, and kinetic ballooning/interchange instabilities (Roux et al., 1991; Voronkov et al., 1997; Cheng, 2004; Saito et al., 2008; Pritchett and Coroniti, 2010) and electromagnetic ion cyclotron

instability (Le Contel et al., 2000; Perraut et al., 2000). Optical measurements from ground-based all-sky cameras (Liang et al., 2008; Rae et al., 2010; Kalmoni et al., 2015; Nishimura et al., 2016) have shown that the optical wave properties (period between 18 and 23 s, azimuthal wavelengths between ~ 60 and ~ 100 km, growth rates $\sim 0.04\text{ s}^{-1}$ and duration of 1 to 1.5 min) are in best agreement with the kinetic ballooning instability (Pritchett and Coroniti, 2010, 2011, 2013). Kinetic instabilities are likely to play a role since the optical wavelengths map to cross-tail distances comparable to the $\sim 2,000$ km ion gyro-radius at $8 R_E$, which is inside the region in the Earth’s magnetotail between ~ 6 and $10 R_E$, where the Earth’s magnetic field often transitions between a quasi-dipolar geometry to a tail-like one. It is in the neighborhood of this region where the onset of substorms is widely acknowledged to occur (e.g., Petrukovich and Yahnin, 2006 and references therein). Models’ predictions for the cross-tail wavelength, growth rate, and frequency of instabilities associated with bead development can be properly tested with an array of spacecraft in the nightside magnetosphere with the following attributes: (1) A main spacecraft emitting an electron beam to ensure the unambiguous magnetic-field connection between the beads observed with ground-based auroral cameras and the magnetospheric region where the beam was emitted from; (2) At least two daughter spacecraft separated azimuthally $\sim 350\text{--}800$ km to allow sufficient resolution to measure azimuthal variation of plasma density, pressure, and convection over the observed $\sim 1,250\text{--}3,200$ km range for beads’ wavelength projected to the magnetospheric equator (Rae et al., 2010; Kalmoni et al., 2015; Nishimura et al., 2016). Similar variable radial separations among spacecraft will allow the necessary resolution of radial gradients. Field and plasma measurements with a 30 s cadence are desirable to test wave growth rates of ballooning and CFCI instabilities in the inner plasma sheet (Rae et al., 2010; Kalmoni et al., 2015; Nishimura et al., 2016).

Auroral and *in-situ* measurements in the magnetosphere suggest that a local decrease in entropy may in some instances influence the triggering of instabilities that cause the onset arcs (**Figure 4**). Observations indicate that a large fraction ($\sim 84\%$) of onsets are preceded by equatorward moving auroral forms (streamers; Nishimura et al., 2010; Xing et al., 2010; Lyons et al., 2011). The decrease in entropy may be caused by low-entropy flux tubes that are injected from the far tail reconnection and intrude into the transition region. In this framework, the streamers observed in the auroral region are assumed to be the low-altitude projection of the low-entropy flux tubes that are moving in the magnetosphere from their source, in reconnection sites farther than $10 R_E$ in the magnetotail, toward the near-Earth environment. The streamers show total field-aligned currents of a few tenths of MA, thicknesses of $\sim 100\text{--}600$ km, field-aligned current densities ranging from less than $1\text{ }\mu\text{A/m}^2$ to more than $20\text{ }\mu\text{A/m}^2$, and a potential drop of a few kV across the stream (Amm et al., 1999; Sergeev et al., 2004). Models of low-entropy earthward propagating flux tubes, sometimes termed “bubbles,” have produced similar Region-1/Region-2 current systems (e.g., Yang et al., 2012). Models’ predictions

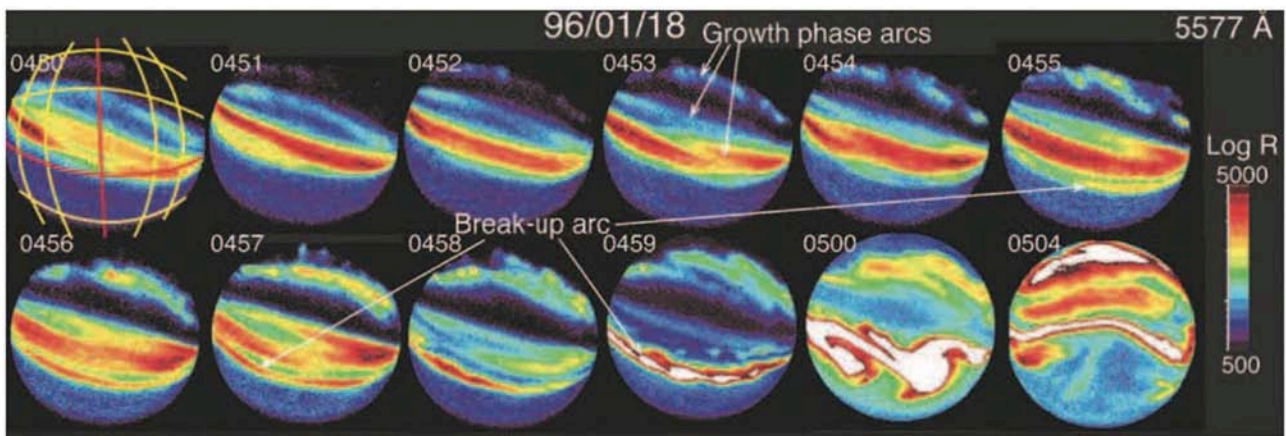


FIGURE 1 | Auroral arcs observed with all-sky imagers in the 557.7 nm band between the last few minutes of growth phase of substorms and the first few minutes of expansion phase (Reproduced from Lyons et al., 2002). Emission is shown in units of Rayleighs.

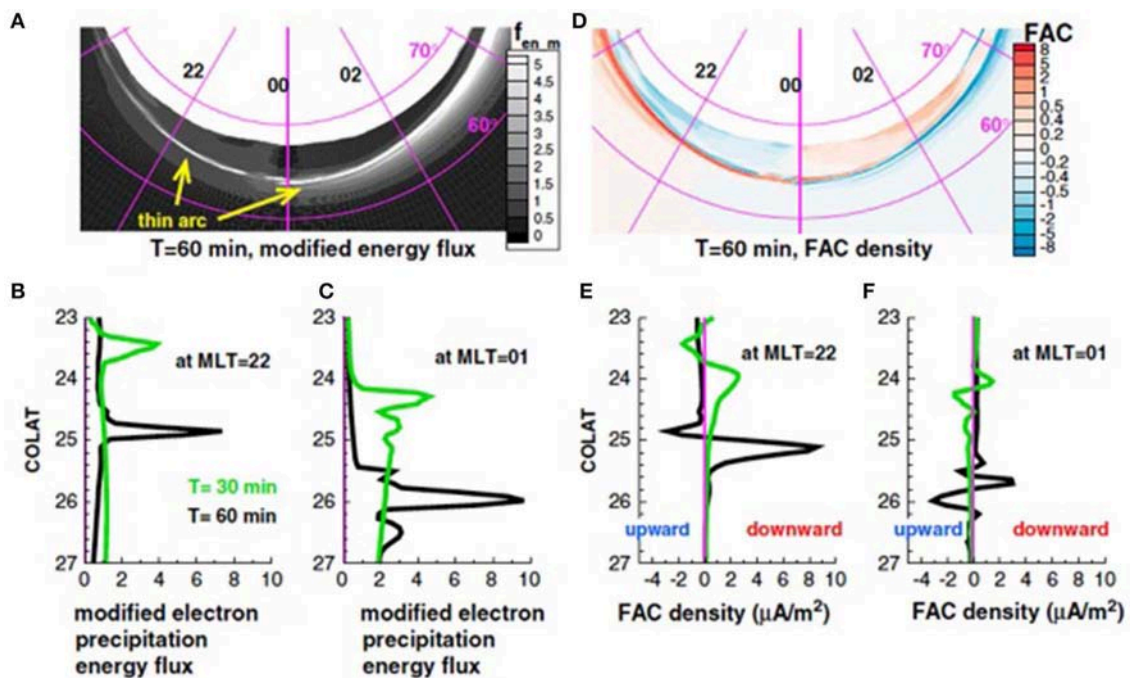


FIGURE 2 | Ionospheric view of (A) modified precipitating energy flux in units of $\text{erg}/\text{cm}^2/\text{s}$, (D) FAC density toward the end of growth phase in units of $\mu\text{A}/\text{m}^2$ (blue is out of the ionosphere). Latitudinal slices of (B,C) energy flux and (E,F) FAC density at MLT = 22 and MLT = 1 at two different times (Reproduced from Yang et al., 2013). The vertical axis COLAT denotes degrees away from the Earth's magnetic dipole axis.

for the rate of change of Hall and Pedersen conductance, horizontal currents, and FACs as well as for the energy flux of precipitating electrons produced by the arrival of bubbles into the inner magnetosphere would be properly tested with the array of spacecraft in the transition region of the nightside magnetosphere, measurements of convection, conductance and FACs in the ionosphere and, most importantly, a method to ensure that the region measured in the magnetosphere maps unambiguously to the region measured in the ionosphere.

The azimuthal separation prescribed for testing the predictions of instability models would be sufficient to measure pressure gradients of $\sim 1.5\text{--}2.5\text{ nPa}/\text{km}$ invoked in the entropy decrease models of onset arc. The same azimuthal spacecraft separation would also resolve the cross-tail structure of the incoming low-entropy flux tubes and, through the electron beam mapping, determine if the low-entropy flow channels measured *in-situ* actually correspond to auroral streamers observed by ground-based imager networks.

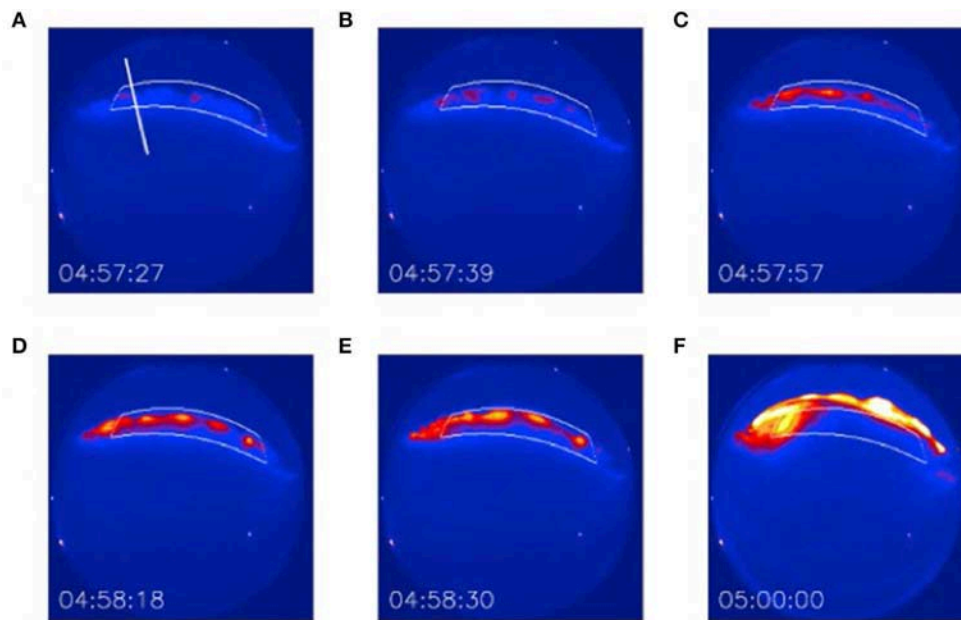


FIGURE 3 | Auroral beads observed by an all-sky imager along an onset arc on 2 October, 2011 in arbitrary units of luminosity in the 557.7 nm emission band. Lines of geomagnetic latitude at 67.8° and 68.4° and geomagnetic longitude at -33.0° and -24.0° define the field of view of the images and show that the onset arc is aligned with constant geomagnetic latitude. The formation and evolution of the beads are observed with time (A–F). The white line in A shows the orientation perpendicular to the arc. After 04:58:30 UT (E) the aurora expands poleward out the box, as can be seen at a later time (F; Reproduced from Kalmoni et al., 2015).

The Substorm Current Wedge

Another crucial outstanding question of substorm development that will be answered with an unambiguous mapping between the magnetosphere and the ionosphere is how the near-Earth M-I coupled system evolves toward a large-scale SCW (McPherron et al., 1973). The SCW is part of a magnetosphere-ionosphere current system that forms during substorm expansion and comprises a current from space into the ionosphere at the eastern edge and out from the ionosphere into space at the western edge of the aurora.

Multiple MHD modeling efforts have related the transport of mass and magnetic flux from the tail to the near-Earth to explain the formation of the SCW (e.g., Birn and Hesse, 1991, 1996, 2005, 2013, 2014; Birn et al., 1999). However, the exact relationship between tail reconnection and near-Earth breakup onset remains to be elucidated. Models show that the SCW configuration that starts with a canonical current polarity (into the ionosphere in the eastern edge of the aurora, out of the ionosphere in the western edge) develops finer structure over the span of a few minutes. The actual SCW current system may actually contain more circuit elements than the standard traditional picture, because of the combined effect of dipolarization, azimuthal flow diversion, shear flows, and twisted/sheared magnetic field. Dipolarization is the process where magnetic geometry changes from tail-like to dipole-like as earthward convection transports magnetic flux from a reconnection site in the tail to near-Earth (Figure 5A). The magnetic shear between the dipole-like geometry inside the SCW and the tail-like geometry outside generates currents that flow into the ionosphere on the eastward edge of the SCW and out of the ionosphere on the westward edge, i.e.,

Region-1 polarity. As the flow is transported closer to Earth it gets diverted to the flanks by the increased magnetic pressure of the ambient dipole magnetic field resulting in rotation of the magnetic field away from the local meridian plane (Figure 5B) and field line twisting generated by earthward and azimuthal flow at the edges of the SCW (Figure 5C). The combined motion generates pairs of oppositely oriented field-aligned currents. The combination of all these effects produce a composite current system as shown in Figure 5D. The outermost (red) current system 1 represents the perturbed currents of the traditional SCW which flows into the ionosphere on the dawn side, flows westward in the ionosphere and flows back to the magnetosphere on the dusk side. The current system 2 (green) is formed by a diversion of radial perpendicular currents into a pair of currents with opposite polarity and closes in the north-south direction in the ionosphere. Current system 3 (blue) is a dusk-to-dawn current near the equatorial plane which is a consequence of the tailward propagation of the dipolarizing region and the associated reduction of the tail-aligned component of the magnetic field. Current system 4 (black) is confined to the equatorial plane and opens the possibility that the current system 2 may close azimuthally in the magnetosphere through current system 4 rather than radially. Models' predictions for the different FAC circuits would be properly tested with an array of spacecraft in the nightside magnetosphere with the following attributes: (1) An appropriate separation to discern particle pressure and flow gradients on the magnetospheric equator as well as spatio-temporal deformation of the magnetic field; (2) An electron beam that ensures the unambiguous identification of the ionospheric foot-point of the regions measured in the

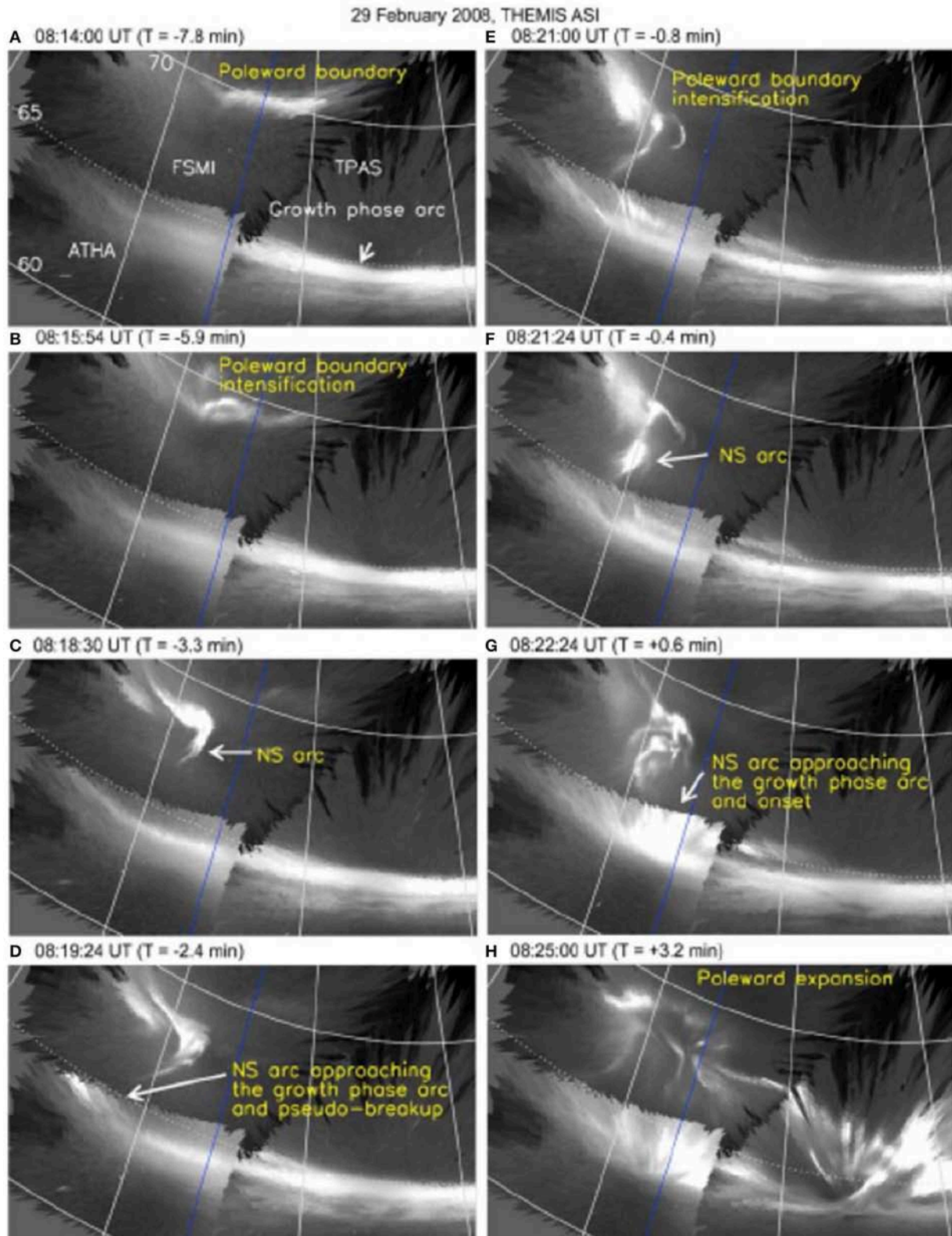


FIGURE 4 | Auroral streamers advancing southward from their origin, the polar cap boundary, to the locus of substorm onset, the growth phase arc (Reproduced from Nishimura et al., 2010). Auroral streamers advancing southward from their origin, the polar cap boundary, to the locus of substorm onset, the growth phase arc (A–G). Onset occurred at 0821:48 UT and expansion is apparent at 0825:00 UT (H) (Reproduced from Nishimura et al., 2010).

magnetosphere; and (3) Appropriate measurements FACs at and around the ionospheric foot-point of the different regions of the SCW. Azimuthally and radially separated measurements of particles and fields enabled by the spacecraft arrangement used for the substorm instability triggering problem would allow the calculation of magnetic-field-aligned currents generated by divergence of current in the magnetosphere and provide *in-situ* measurements of vorticity and pressure gradients that contribute to the complex system of currents in the substorm current wedge.

ELECTRONS AS PROBES OF THE MAGNETIC FIELD

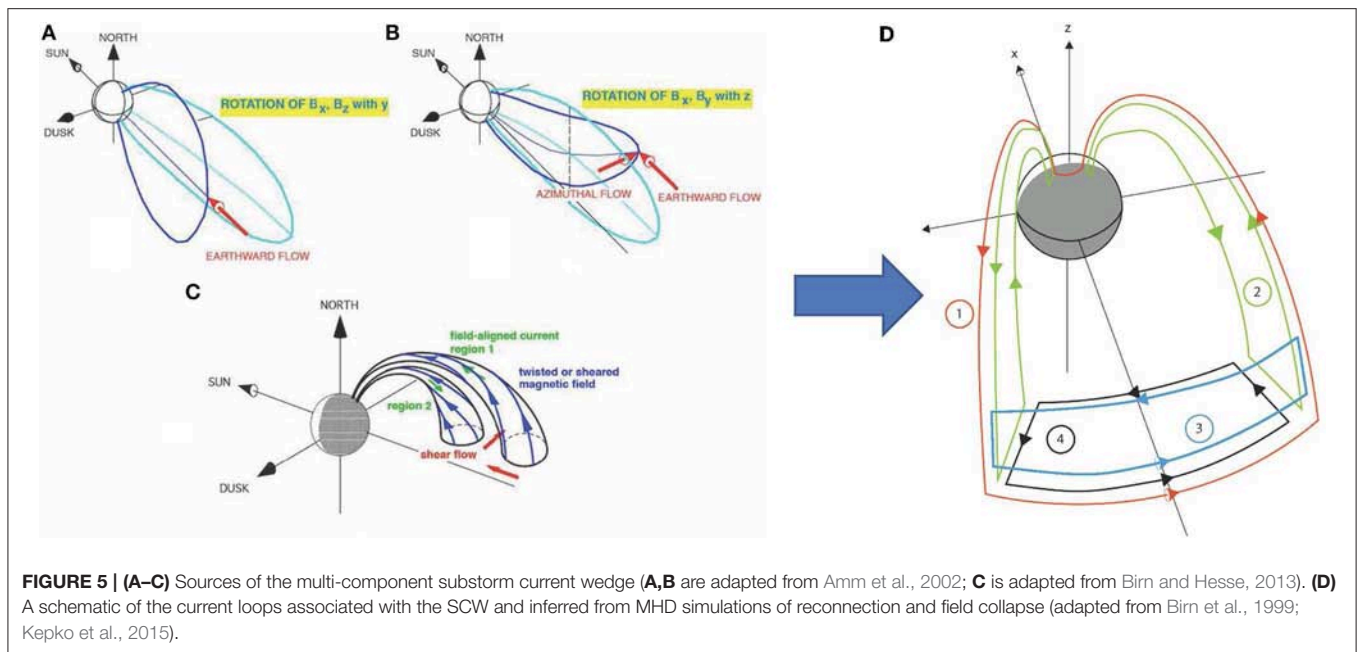
All the comparisons between observations and predictions for substorm growth, onset and current wedge theories have involved space and time histories of optical auroral features coupled with *ad-hoc* mappings between *in-situ* measurements in the magnetotail and the observed auroral forms. Various empirical and magnetohydrodynamic (MHD) modeling techniques have been used to approximate the instantaneous configuration of the magnetospheric magnetic field (e.g., Tsyganenko, 1989; Sergeev et al., 1993; Fedder et al., 1995; Kubyshkina et al., 1999; De Zeeuw et al., 2004; Toffoletto et al., 2004; Pembroke et al., 2012). Despite advances in field mapping, the uncertainties involved are still quite large. Difficulties arise due to dynamic phenomena, especially in the tail region, in the form of thin current sheets, magnetic flux ropes, non-adiabatic substorm and storm magnetic field reconfigurations, and high-speed flows (e.g., Donovan et al., 1992; Jordan et al., 1992; Peredo et al., 1993; Fairfield et al., 1994; Pulkkinen and Tsyganenko, 1996; Thomsen et al., 1996). At geosynchronous altitude, the statistical uncertainty in the mapping given by magnetic field models is $\sim \pm 3^\circ$ (e.g., Reeves et al., 1996). Up to 20% of the time-field models could be off by more than 5° . Non-adiabatic conditions also mean field lines will no longer be equipotential, making it hard to causally relate magnetospheric and ionospheric flows (e.g., Hesse et al., 1997). Recent efforts to couple models of ion drift physics in the inner magnetosphere to MHD models of the outer magnetosphere (De Zeeuw et al., 2004; Pembroke et al., 2012) reveal new structure and dynamics in the magnetotail; the same efforts underline inherent complexity in the magnetic-field topology. This communication describes how unambiguous magnetic field mapping can be achieved by firing a beam of high-energy electrons from the source region into the ionosphere. It also describes how the properties of the magnetosphere source region can be properly described by deploying a constellation of spacecraft in the vicinity of the spacecraft that fires the beam.

Multiple active experiments that include the injection of artificial energetic electron beams from sounding rockets to investigate magnetospheric structure and dynamics were conducted by a number of groups in the 1960s and 1970s (see Winckler, 1980, for a review). These experiments used keV electron beams and focused on tracing magnetic-field

lines by injecting and detecting mirrored electrons; using beams as diagnostic tools to sense local electric and magnetic fields; and investigating wave-particle interactions, including the generation of electromagnetic radiation, the scattering of energetic electrons by waves, and general beam-plasma physics. The optical signature of electrons emitted from a rocket in the polar region of one hemisphere was detected in the conjugate ionosphere and in the hemisphere of origin after bouncing along field lines (Hallinan et al., 1978, 1990). Rocket experiments with electron beams, emitted at energies up to several-keV and currents up to 1–2 Amperes, were carried out in the 1980s and 1990s to explore the effects of these beams on the vehicle near-plasma environment and the upper and middle ionosphere (e.g., Mandell et al., 1990; Neubert et al., 1991; Neubert and Banks, 1992; Raitt et al., 1995). Shuttle/Spacelab electron beam experiments with similar energies and currents were carried out in the same time period to further measure and model the effects of beams (Neubert et al., 1986, 1995; Bush et al., 1987; Cai et al., 1987; Reeves et al., 1988, 1990; Burch et al., 1993). Beam effects studied under these experiments included beam-induced space charge, generation of artificial aurora, and generation of VLF waves by pulsed beams. In the late 1990s and early 2000s some theoretical studies considered the applications and technical challenges associated with injection of relativistic electron beams into the space environment (e.g., Neubert et al., 1996; Krause, 1998; Krause et al., 1999; Gilchrist et al., 2001; Neubert and Gilchrist, 2002, 2004). Results from these studies indicate that relativistic beams should be far more stable than keV beams due to a combination of the higher relativistic electron mass, lower beam densities, and less pronounced spacecraft-charging effects, at least for injections from the ionosphere.

Compact linear accelerators are currently capable of generating beams with currents ($\lesssim 100$ mA), energies (1–10 MeV), pulse durations (μ s), and duty cycles ($\sim 0.1\%$) that make them the best candidates for application to magnetic field mapping. Because of increased efficiency, high frequency, and high gradient technologies developed since the 1990s (Ruth et al., 1993; Wang, 2009; Dolgashev and Tantawi, 2010; Neilson et al., 2010; Tantawi and Neilson, 2012), compact linear accelerators can fit in a Mid Explorer class mission's size, mass, and power envelope. However, the realization of the proposition that beams of relativistic electrons can be used as magnetic field tracers require demonstration that the beam can be injected into the loss cone from magnetospheric altitudes, that the spacecraft potentials induced by the beam emission are manageable, and that sufficient electron flux reaches the atmosphere to be detectable by optical or radio means after the beam has propagated thousands of kilometers under competing effects of beam spread and constriction as well as effects of beam-induced instabilities.

In the next section we provide a review of the latest results of synergistic research carried out under the NSF INSPIRE program to address these challenges and discuss the next steps toward the realization of active experiments in space using relativistic electron beams.



CHALLENGES WITH THE INJECTION OF SPACE-BASED BEAMS

Plasma Response Time and Spacecraft Charge Control

Past electron beam experiments encountered issues with the injection of space-based beams due to rapid spacecraft charging, which influences beam fidelity and beam-plasma interactions (e.g., Cohen et al., 1980a,b; Gussenhoven et al., 1980). These issues led to studies of the fundamental processes associated with electron-beam-plasma interactions, including the formation of sheath regions of particle and field fluctuations, plasma-neutral gas interactions, wave-particle interactions, and non-linear phenomena (Neubert and Banks, 1992). Some of these issues are associated with charge and current neutralization: in a plasma, the axial field of an electron beam can effectively expel thermal electrons to become charge-neutralized (Humphries, 1990). SCATHA spacecraft experiments carried out in 1979 to investigate the effect of the interaction between the magnetospheric plasma and keV beam emission on spacecraft potential demonstrated that, for certain beam currents, the plasma can supply the return current required to keep the spacecraft potential below the beam energy so the beam can safely be emitted (Cohen et al., 1980a), but for higher beam currents most of the beam electrons return to the spacecraft (Gussenhoven et al., 1980) or even cause the failure of the spacecraft systems (Cohen et al., 1980b).

The ability of the plasma to respond to an injected beam of electrons depends on the plasma response time, which is driven by the plasma frequency (Humphries, 1990). Charging for single-pulse injections is expected to have a negligible effect on beam fidelity due to the low charge-accumulation build up ($\sim$ kV spacecraft potential) compared to the beam energy ($\sim$ MeV). The charging resulting from extended pulse emissions depends on the

ability of the ambient plasma to supply the return current, given by the thermal current density. Neubert and Gilchrist (2002) have shown that spacecraft charging and beam-plasma interactions become significant for currents of ~ 100 A for ionospheric injections of 5 MeV beams, far larger than the expected electric current demands of space-based electron accelerators.

For injections from the magnetosphere, the effects of spacecraft charging on beam fidelity must be considered, but with flexible beam operations (beam energy, duty cycle, etc.), paradigms for stable injection are expected. Recent studies by Delzanno et al. (2015a,b) have demonstrated an operational paradigm where releasing a high-density neutral contactor plasma prior and during beam ejection leads to successful beam ejection. As such, several-MeV beams should be suitable for injections from the ionosphere and the magnetosphere.

Beam Injection Into the Loss Cone

To maximize the fraction of the relativistic electron beam entering the atmosphere, the beam must be injected into a geometrical region known as the loss cone. Beams injected outside the loss cone will bounce back to their source, due to the magnetic mirror force, before they reach the atmosphere. Standard calculation of the loss cone involves the conservation of the first adiabatic invariant (e.g., Rossi and Olbert, 1970). For sub-relativistic particles the calculation of the width of the loss cone is sufficiently accurate using the zeroth order term, $\mu^{(0)}$, of the adiabatic invariant corresponding to the cyclotron motion, μ , which is an asymptotic series in the small parameter ρ/L_B (Northrop, 1963) where ρ is the effective Larmor radius, defined by v/Ω_c , where v is the initial velocity of the particle (total, not just $v_{\perp}$), Ω_c is the cyclotron frequency of the particle $|q_{\parallel}B|/mc$, q is the particle's charge; and m is its relativistic mass, given by $m_0\gamma$, where m_0 is the rest mass of the particle, and $\gamma = 1/\sqrt{1 - v^2/c^2}$. The denominator, L_B , of the small parameter is the characteristic

gradient length scale of the magnetic field $L_B^{-1} = |\nabla \ln B|$. For relativistic electrons Porazik et al. (2014) showed that higher-order terms of the magnetic moment invariant are necessary to correctly determine the mirror point of trapped energetic particles, and therefore the loss cone. **Figure 6** (left) shows the pitch angles (δ) that would lead to precipitation for different azimuthal injection angles (λ) as a function of electron energy at $10 R_E$. The electron is considered to be lost if its mirror point is at the radial distance of $1 R_E$ or less. For comparison, the dashed line shows the loss cone computed based on only the lowest-order term of the magnetic moment. The importance of higher-order terms is most dramatically reflected in the λ -dependence of the loss cone. As the energy of the electron increases, the λ -dependence becomes more pronounced, and eventually the loss cone becomes a closed contour with unique boundaries in both angles. The largest range of δ angles is always at $\lambda = -90^\circ$, in the direction of electron drift, tangentially to the flux surface (for positive ions, the optimal value of λ would be $+90^\circ$). **Figure 6** (right) shows the loss cone for a 7-MeV electron initialized in the equatorial plane at different distances. As the distance increases the loss cone again becomes confined to a small region in phase space, with unique boundaries in both angles. The modified loss cone resulting from the inclusion of higher order terms is no longer entirely defined by the traditional pitch angle but also by the phase angle of the particle at the point of injection. The optimal orientation of the injection has a non-zero component perpendicular to the magnetic field line, and is in the plane tangential to the flux surface. The results show that injection-angle control is important and depends on location and beam energy. Also, as we will discuss in the next section, this theory is backed up by single-particle simulations, which do not rely on the assumption of conservation of the first adiabatic invariant.

The results of these studies are important in guiding the design considerations that determine the energy and pointing envelopes that ensure electrons trajectories' reaching the atmosphere when injection occurs in the magnetosphere. The fraction of the beam reaching the atmosphere will depend on the fraction of its phase space density lying inside the modified loss cone at injection. For instance, **Figure 6** shows that a 7 MeV beam injected at -90° azimuth at $10 R_E$ must have a spread smaller than 2.6° to ensure precipitation into the atmosphere. Current accelerators can achieve spreads that are approximately an order of magnitude smaller than 2.6° thus enabling the entire particle flux to be inside the loss cone.

Using the physical parameters of a reference 1 MeV electron beam instrument, Powis et al. (2019) determine that an appropriate treatment for injection of relativistic particles in a dipolar field at the geomagnetic equator must include the first three terms in the expansion series for μ in order to capture small changes in δ and λ (Powis et al., 2019, Equation 12). As particles approach Earth, the most significant contribution to the magnetic moment is given by the zeroth order component because of the combined effect of a stronger magnetic field and a larger the perpendicular velocity. Relating the initial magnetic moment, $\mu^{(0)} + \mu^{(1)} + \mu^{(2)}$, to the final magnetic moment, $\mu_{\oplus}^{(0)}$, gives a general relationship for the final cyclotron radius at impact for any particle injected from the equatorial plane

along a dipole field line, as a function of injection radial distance from Earth, injection energy and injection angles λ and δ . For a particle injected directly along the field line $\delta = 0$ and reference properties, 5.7% of the initial parallel kinetic energy is converted into perpendicular kinetic energy at impact, resulting in a cyclotron radius, r_c , of 21.8 m. Increasing the value of δ results in an increase of the beam radius at Earth's impact. Powis et al.'s calculations show, for instance, that increasing δ to 1° results in a final cyclotron radius of 60 m.

Beam Propagation

The electron beam produced by a radio-frequency (RF) linear accelerator is a concatenation of periodic pulses the smallest of which are picosecond length micro-pulses. Multiple micro-pulses are bunched together to form a mini-pulse, typically lasting several microseconds, and a group of mini-pulses constitute a macro-pulse. Multiple combinations of macro-pulse duration and repetition rate can be chosen according to science objectives, from sub-second to multiple seconds. To illustrate the propagation properties of the beam and to characterize its properties at the point where it comes in contact with the Earth's topside atmosphere we have used a set of illustrative beam parameters described in Powis et al. (2019).

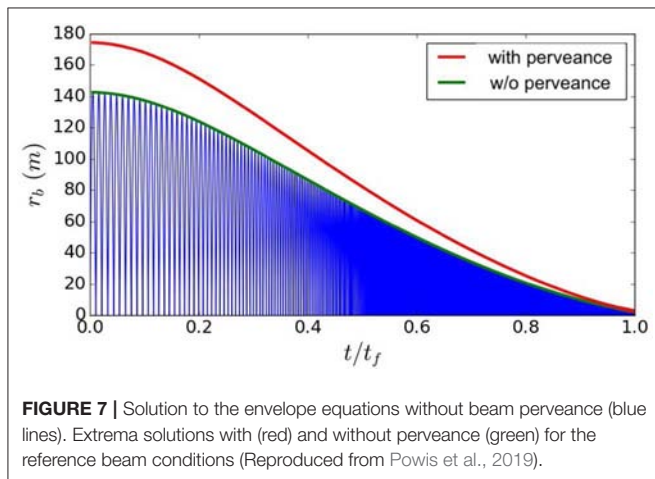
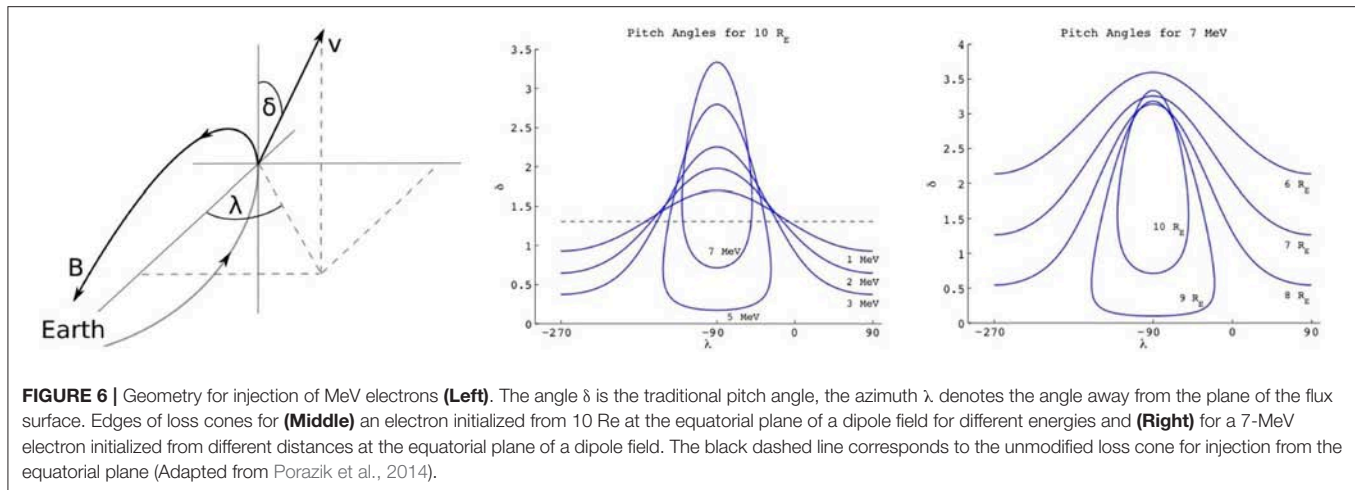
One approach that we have adopted to describe the beam propagation assumes that the evolution of the mini-pulse's RMS radius, r_b , can be decoupled from the electrons' helical motion so, within the frame of the beam centroid the evolution of the mini-pulse distribution depends on the initial conditions, the self-generated electromagnetic forces, and the ambient magnetic field. Since the mini-pulse length remains much longer ($\sim 10^3$) than the beam's radius even at maximum expansion, the mini-pulse is modeled as an infinitely long beam. Under these circumstances the radius r_b can be considered as the envelope of the beam and the standard one-dimensional beam envelope equation is applied (Reiser, 2008),

$$\frac{d^2 r_b}{dS^2} = -k_0^2 r_b + \frac{K}{r_b} + \frac{\epsilon_r^2}{r_b^3} \quad (1)$$

Where S is the arc length from its injection point to its current position, $\epsilon_r = \frac{v_{\perp} r_b}{v_0}$ is the beam radial emittance, $v_0 = \beta c$, $k_0 = \frac{qB(S)}{2mc\gamma\beta}$ parametrizes the focusing due to the ambient magnetic field and $K = \frac{qI_0}{2\pi\epsilon_0 m_0 c^3 \beta^3 \gamma^3}$ is the perveance, which captures the influence of beam self-charge and self-magnetic field (Powis et al., 2019, and references therein).

The solution to Equation (1) produces an oscillating beam envelope that initially grows to a size greater than a hundred kilometers but progressively narrows as the beam propagates into stronger magnetic field when it approaches Earth (**Figure 7**). The final beam radius at Earth's topside atmosphere impact is 2.6 m.

Equation (1) is also solved for the case without perveance ($K = 0$) to allow for comparison with ballistic simulations resulting from a single-particle propagation algorithm which does not incorporate the effects of space charge. The comparison is used to determine the appropriateness of ballistic simulations to describe the beam propagation. Parametric solutions for the envelope



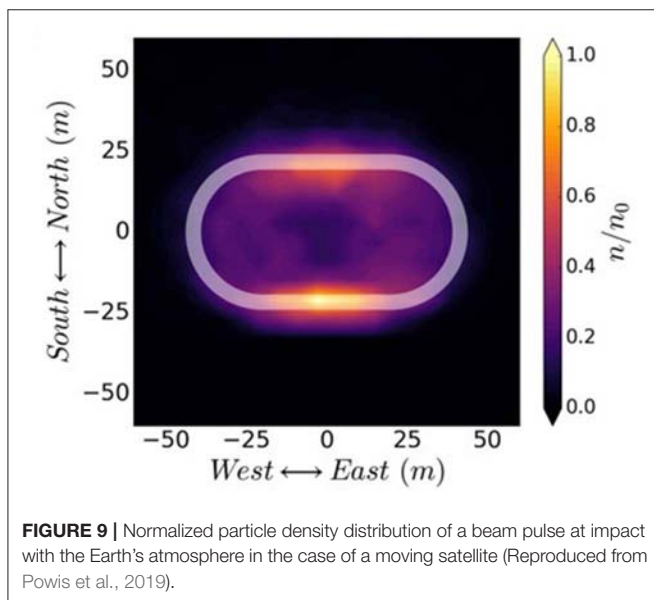
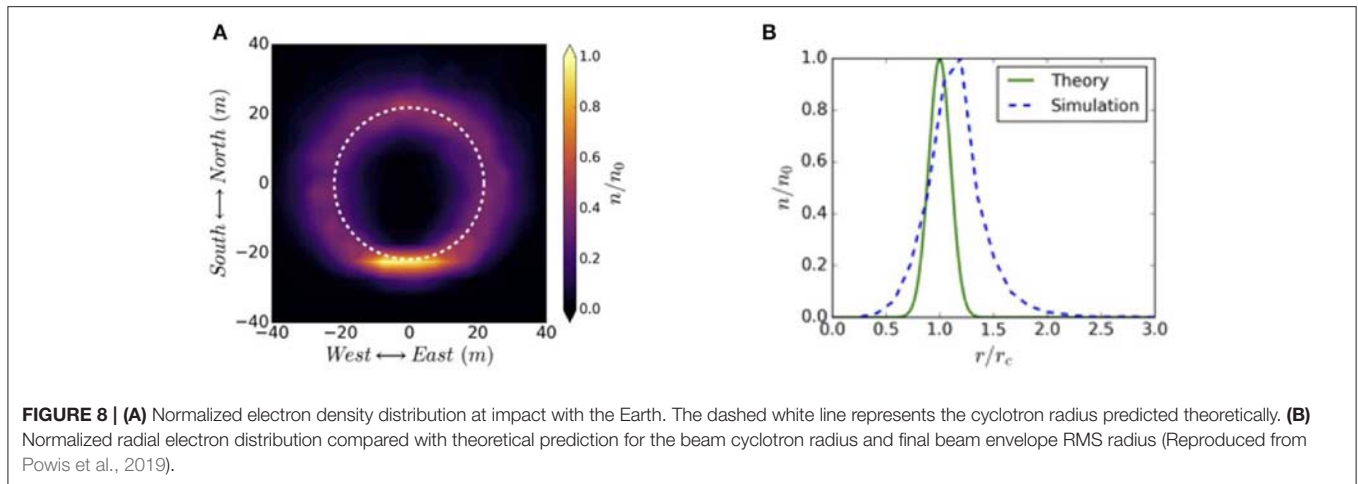
equation show that for the reference beam current of 1 mA, the final beam envelope radius is only weakly affected by beam perveance. This is due to the small magnitude of the average current and to the fact that, for relativistic beams, the self-generated magnetic field acts to cancel out a large fraction of the beam self-charge. This also demonstrates that despite self-forces being neglected, the use of single-particle simulations is suitable for modeling beams with similar properties to those of the reference values. Two other important properties are found. Firstly, increasing the initial beam energy results in a larger final radius since the increased electron momentum reduces the effectiveness of the applied magnetic field to focus the beam. Secondly, increasing beam current results in an increased final radius due to the increased current density at the point of emission. The larger the beam current, the less suitable ballistic simulations are for modeling the beam.

The properties of the cross-section density profile expected for a beam arriving to the topside atmosphere can be explored by applying the single particle algorithm to an ensemble of electrons from a mini-pulse injected from 10 R_E along a dipole field line using the physical parameters of the reference 1 MeV electron

beam instrument described by Powis et al. (2019). The resulting density distribution (shown in **Figure 8A**) is ring-shaped rather than circular. The white dashed line shows the cyclotron radius $r_c = 21.8$ m expected for a 1 MeV electron at the topside ionosphere after having propagated along a dipolar field with a conserved first adiabatic invariant. The density distribution given by $n/n_0 = \exp\left((r - r_c)^2 / 2r_{bf}^2\right)$ (shown in **Figure 8B**) suggests that the beam radius at Earth's impact obtained from the simulation is 15% larger than 21.8 m. Cyclotron radius calculated at the topside ionosphere is expected to be the dominant parameter in determining the final beam spot size at the top of the atmosphere. There is less clear agreement between the simulated and predicted RMS beam envelope radius, as the simulation shows a larger envelope radius than the 2.2 m predicted by the envelope equation. Such discrepancy may be due to the effect of energy dependence of the ∇B and curvature drifts. Since the beam particles have an initial spread in energy, that spread translates into a smearing of the beam's final density distribution. The reason for the beam's ring shape can be explained by a combination of the following effects. Initially the envelope of the ensemble of electrons is expanding and contracting with the cyclotron frequency and in phase with the centroid motion of the entire beam rotating at r_c . Note that for the given beam parameters in the simulation, $r_c > r_b$. Since there is energy spread in the beam the cyclotron frequencies of particles will also have a spread due to small variations in the γ factor. As the beam propagates along the magnetic field line the particles will spread many periods in gyro-phase resulting in a beam density profile that evolves into a ring.

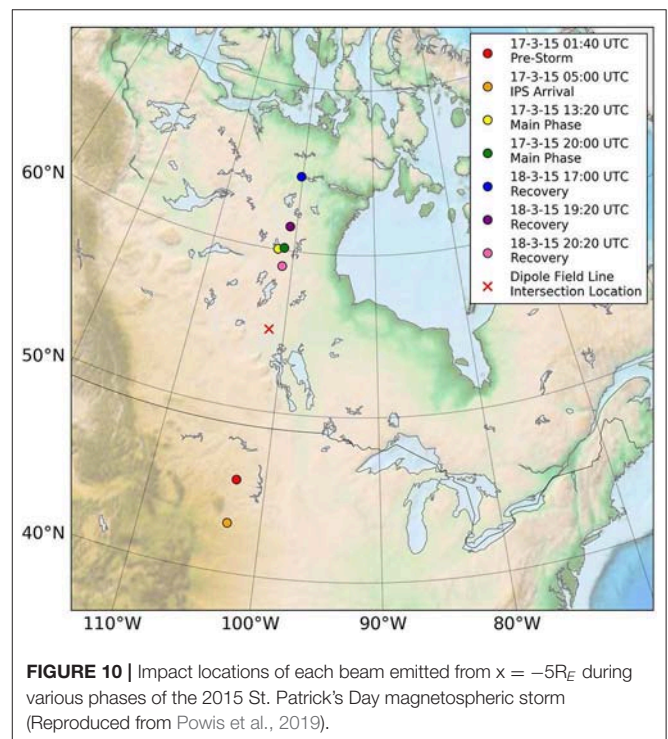
We note that in **Figure 8A** density appears preferentially concentrated on a spot at the bottom of the ring when beam self-forces are not included. This occurs despite the expected decorrelation of particle gyro-orbit phase because a large number of particles remain closely correlated after traveling the length of the field line. Including self-forces (perveance) will in reality generate a more uniform density distribution around the ring.

An additional property of the cross-section density profile that must be considered for beams emitted from the magnetosphere



into the atmosphere is the east-west spread of the beam caused by the spacecraft motion relative to the Earth's atmosphere. Simple Newtonian mechanics calculations show that the east-west elongation of the density distribution increases slightly, to ~ 72 m, for a 0.5-s burst of mini-pulses from a spacecraft orbiting at $10 R_E$ altitude (**Figure 9**). Other less pronounced effects, such as the gradient and curvature particle drifts, contribute to the elongation of the distribution as well.

The beam propagation calculations done in a dipolar field can be extended to a more realistic magnetic fields that are generated self-consistently by global MHD simulations such as BATS-R-US (Tóth, 2005). Beam injections from $5 R_E$ were simulated for varying magnetic field configurations experienced at different stages in the development of a geomagnetic storm in March 2015. The simulation provides a picture of the range of variation in the latitude of atmospheric foot-point of the beam as well as the size of the beam. The spread of $\sim 10^\circ$ on either side of the foot-point of geosynchronous altitude (**Figure 10**) shows the large



variation in the geographical location of the beam's signature due to the significant changes in magnetic field topology induced by geomagnetic activity. Prediction of beam foot-point location as a function of geomagnetic activity level provides guidance on where ground-based imagers should be placed to ensure that the beam's optical signature would be captured.

Beam Detectability and Ground-Based Diagnostics

Neubert et al. (1996), Krause (1998), and Krause et al. (1999) demonstrated that relativistic beams injected from the ionosphere into the atmosphere below would produce

significant electron-density enhancements, optical emissions, and measureable height-integrated X-ray fluxes.

Marshall et al. (2014) expanded on the work of Krause (1998) to calculate optical emissions observable from the ground, X-ray production and propagation and detectability from satellites and balloons, and backscattered electrons that could be observed from Low Earth Orbit (LEO). That study showed that optical signatures were likely detectable, X-ray fluxes were likely to be far too low from either LEO or balloon altitudes, and ionization could likely be measured from the ground using incoherent scatter radar. That study investigated a pulse of electrons with 0.05–1 Joules of total energy. Recent accelerator design efforts are targeting a beam total energy of 100–1,000 J, prompting a revisit to the calculations of Marshall et al. (2014).

In the new simulations (Marshall et al., this issue), the accelerator under consideration produces an output of 5 J of electrons at 1 MeV in each pulse (3.1×10^{13} electrons), with a pulse every 5 ms. A beam of 1 MeV electrons injected from a distance of $10 R_e$ was simulated by Porazik et al. (2014), who then propagated ballistically the beam to 300 km altitude and calculated the spatial, energy, and pitch angle distributions of the beam at that altitude. Those distributions are used as the input distributions to Marshall et al.'s Monte Carlo modeling. A 2-D histogram of the particle positions shows that the beam is distributed approximately as a Gaussian with a 1-sigma diameter of 311 m at 300 km altitude. The beam is extremely field-aligned, with a divergence of less than 1 degree, due to the careful choice of the firing direction in Porazik et al. (2014). However, simulations show that as long as the beam is inside the loss cone, the pitch angle distribution plays only a small role in the atmospheric signatures. For example, a beam with all electrons 60-degree pitch angle at 300 km altitude, just inside the loss cone, will have a similar energy deposition profile, but raised in altitude by 4 km.

The new simulation results show that the peak of the energy deposition from a sequence of 20 pulses spanning 100 ms and totaling 100 J, or a sequence of 200 pulses spanning 1 second and totaling 1 kJ, occurs slightly below 60 km altitude, in the atmospheric region known as the D-region, and that approximately 2.2% of the total injected energy is converted to N₂ 1P emissions, and 0.6% is converted to N⁺2 1N emissions. Using these parameters and considering an optical aperture of 50 mm diameter (a typical camera lens) with a field-of-view that is larger than the emitting region, and where one can expect 3.3×10^3 photons to be collected by the instrument, Marshall et al. conclude that a PMT-based system can detect the emission produced by the beam with a signal-to-noise-ratio (SNR) of 25 when sampled at 100 Hz. For a 50 mm diameter lens wide field-of-view camera system an SNR of 10 is feasible when sampling at 10 Hz.

Marshall et al. also consider whether the electron density enhancement that would be produced by the beam in the D-region is detectable by radar. Their calculations show that the expected electron density enhancement after the 1-s train of pulses and the ~ 1 -s recovery in the D-region for the 100 J (20 pulses in 100 ms) beam emission case. After 20 pulses, the peak electron density of $3.9 \times 10^9 \text{ cm}^{-3}$ occurs at an altitude of 59 km. It also shows the SNR that an incoherent scatter radar (ISR) such

as the Poker Flat ISR (PFISR) would measure and the expected relative error, dS/S . Although $\text{SNR} < 1$, detectability is actually determined after calculating the gain for a Lorentzian radar spectrum done by averaging radar pulses coherently (e.g., Farley, 1969). The relative error dS/S is then found by incoherently averaging all the sets of coherent averages embedded in the interval where the radar sampled the ambient electron density enhancement produced by the electron beam. The relative error for the example shown by Marshall et al. is $dS/S = 0.27$. A value of $dS/S = 1$ indicates that the signal is 1σ above zero SNR; $dS/S = 0.33$ indicates 3σ above zero SNR, and so forth. The analysis thus shows that ISRs operating standard beam codes (whose parameters were applied in the calculation of SNR and dS/S) are capable of detecting the beam pulse sequence of 1 kJ injected over 1 s. New radar beam codes that increase the coherent gain, combined with longer integration times and longer electron beam pulses, will decrease dS/S thus improving detectability. For example, a 50% increase in the number of averaging intervals would decrease dS/S to 0.22.

An important environmental consequence of the beam interaction with the atmosphere addressed by Marshall et al. is the possibility of adverse effects on the atmosphere. Energetic electron precipitation leads to enhancement of odd nitrogen (Rusch et al., 1981) and odd hydrogen (Solomon et al., 1982). These molecules are long-lived and, as they are transported downward into the stratosphere, can affect ozone concentration (e.g., Callis et al., 1991, 1996). Marshall et al. apply the Glukhov-Pasko-Inan (GPI) chemistry model (Glukhov et al., 1992; Lehtinen and Inan, 2009) and the Sodankylä Ion and Neutral Chemistry (SIC) model (Verronen et al., 2005; Turunen et al., 2009) to calculate the density increase in NO_x, HO_x and decrease in ozone due to the precipitation of the electron beam. The SIC model shows an increase in NO_x density of only 0.5% from its background density and an increase in HO_x of 0.4%. The ozone signature is negligible. These results show that active experiments with relativistic electron beams pulsed at short intervals can be used for magnetosphere-ionosphere research without causing significant adverse long-term effects in the atmosphere.

ADDITIONAL APPLICATIONS OF BEAM EXPERIMENTS

Relativistic electron beams have multiple applications beyond the fundamental problems in space physics discussed so far. We briefly discuss two of them: Sprite triggering and beam-induced waves to precipitate radiation belt electrons through resonant pitch angle scattering.

Sprite Triggering

Enhanced conductivity channels above thunderstorm systems can lead to the modification of the atmospheric potential structure. The resulting electric fields may lead to atmospheric breakdown and discharge, known as sprite, especially at high altitudes, where the breakdown fields, E_k , are less than 100 mV/m (Banks et al., 1987, 1990; Neubert and Banks, 1992; Neubert et al., 1996). Neubert and Gilchrist (2004) suggested the

possibility that the relativistic electron beam, upon its interaction with the atmosphere, could modify the conductivity enough to enhance the triggering of sprites at their typical triggering altitude of ~ 75 km (Stenbaek-Nielsen et al., 2010; Pasko et al., 2012). Marshall et al. investigate the possibility of triggering sprites with MeV-class beams by calculating the electric fields induced by the beam precipitation above a thunderstorm system using the 2-D quasi-electric (QES) field model of Kabirzadeh et al. (2015, 2017). These results show that after the discharge $E > E_k$ within 1 km of the beam radius which is a condition expected favorable for sprite triggering, thus allowing for the possibility to conduct a carefully timed experiment to increase the high-altitude electric field to trigger sprites.

Wave-Particle Interactions and Loss of Electrons

The radiation belts are near-Earth magnetosphere regions populated by protons and electrons with energies from 100 keV to > 15 MeV. Enhanced radiation-belt electron fluxes, which can be caused by geomagnetic storms or anthropogenic sources, are known to be damaging to space assets (Baker, 2001; Horne, 2003). Particles originating in the solar wind and the ionosphere are accelerated during geomagnetic storms through wave-particle interactions and radial transport and become trapped in the $1.5\text{--}5 R_e$ region (e.g., Horne et al., 2005; Shprits et al., 2008a,b and references therein). Some of these particles can be lost by precipitation into the atmosphere (Lorentzen et al., 2001a; Millan et al., 2002; Green et al., 2004; O'Brien et al., 2004; Bortnik et al., 2006; Millan and Thorne, 2007; Thorne et al., 2010). Theoretical work carried out in the 1960s and 1970s showed that wave-particle interactions can lead to pitch-angle scattering of electrons and their subsequent loss to the atmosphere (Kennel and Petschek, 1966; Thorne and Kennel, 1971). Violation of the first two adiabatic invariants can induce pitch-angle scattering (and potential loss to the atmosphere) and energy diffusion.

Elucidating how wave-particle interactions cause the radiation belts to lose electrons is important for mitigating space weather effects. Considerable research has investigated methods to control radiation belt populations using VLF-wave injection to precipitate these particles (e.g., Inan et al., 1984, 2003). However, challenges exist with methods for efficiently transmitting VLF waves to the space plasma.

Radiation-belt electrons in the $0.1\text{--}10$ MeV range resonate with VLF whistler-mode waves of $0.1\text{--}10$ kHz. The natural environment often contains waves in the VLF band, such as hiss, chorus, and lightning-generated whistlers. The source of hiss and the depletion and refilling rates of the radiation belts are topics of active research. Whistler mode chorus consists of discrete whistler mode emissions observed outside the plasmasphere in the frequency range of $0.1\text{--}1 f_{ce}$ (~ 100 Hz–5 kHz). Models of whistler-electron interaction suggest that whistler mode chorus waves are generated at the equator first, driving the pitch angle scattering of ~ 10 keV electrons, which can cause pulsating aurora (Lessard, 2012). Subsequently, the waves propagate to higher latitudes where pitch angle scattering of sub-relativistic ($\sim$ few hundreds of keV) and relativistic electrons ($\sim$ MeV)

occurs. Whistler mode waves first resonate with electrons at tens of keV near the equator, and then with higher-energy electrons at higher latitudes (Lorentzen et al., 2001a; Horne and Thorne, 2003; Thorne et al., 2005). Therefore, precipitation of electrons across a wide energy range is expected. In most cases, whistler mode chorus is characterized by discrete elements called “risers,” which generally have rising-tone frequency-time spectra between ~ 0.1 and $0.8 f_{ce}$, although falling tones can occur. Outside the plasmasphere, electron resonant energies for typical whistler-mode frequencies near 2 kHz are ~ 100 keV for interactions occurring at the equator. Scattering by whistler-mode chorus was suggested as the mechanism responsible for relativistic electron microbursts, since both are most often observed between 0300 and 1,500 magnetic local time (Lorentzen et al., 2001b).

Controlled electron injections at specified energies and pitch angles would enable detailed studies of wave-particle interactions and scattering. An injected beam of known particle energy and pitch angle can be used to target specific wave frequencies for growth or generation. The use of a modulated (via changes to pitch angle and energy), relativistic electron beam to excite VLF waves may be an efficient method to scatter enhanced radiation-belt electrons into the loss cone. Investigations of the dynamics, stability, and loss of artificially injected relativistic electron beams (Pritchett et al., 1989; Khazanov et al., 1999a,b, 2000) indicate that they could be powerful means for studying wave and collisional interactions. The application of electron beams to trigger waves that can scatter radiation belt electrons into the loss cone is an active area of research (see Delzanno et al., this issue).

CONCLUSION, ROADMAP TO THE APPLICATION OF ELECTRON BEAMS

Active experiments with relativistic electron beams represent the most viable opportunity to finally bring closure to long-standing problems how the magnetosphere and the ionosphere connect to generate aurora, to transfer energy between the two domains, and to regulate the circulation of mass, momentum, and energy throughout the ionosphere-magnetosphere system. A spacecraft mission that measures *in-situ* particle density, pressure, convection, and electric current as well as radial and azimuthal gradients of these quantities with a distributed set of measurements will be able to quantify the source terms that drive the electromagnetic connection with the ionosphere. Accurate correspondence between magnetospheric processes or regions and their ionospheric foot-points can be achieved with beams of energetic electrons emitted in the magnetosphere under controlled conditions, propagating along magnetic-field lines in fractions of a second, and detected by an array of ground-based optical imagers, radars, riometers, or X-ray detectors through the optical, radio, and electron density imprints created in the atmosphere by the impact between the beam's electrons and the neutral particles in the atmosphere. Given the current state of compact-accelerator technology, development, and launch of a space-based energetic particle accelerator are only a few years away. The technology for relativistic linear electron accelerators will overcome the challenges encountered with previous efforts

to trace magnetic-field lines with lower energy electron beams emitted from the magnetosphere.

The research results presented here demonstrate the feasibility of using relativistic electron beams that: (1). When emitted under appropriate conditions do not raise the spacecraft potential more than a few kV; (2). Can propagate along realistic field lines, when emitted inside a modified loss cone geometry applicable only to relativistic particles; (3). Can propagate into the topside ionosphere with sufficient flux to generate a perturbation in the middle atmosphere that is detectable on the ground with optical and radio instruments and such that does not produce a significant lasting adverse effect on the chemistry of the atmosphere.

These results are highly encouraging, and work continues to definitively demonstrate the validity of the beam emission concept as a viable active experiment tool for magnetosphere-ionosphere research applications. One of the areas of ongoing research is the determination of the stability properties of the beam as it travels along magnetic field lines from the region around 6–10 R_e in the night-side magnetosphere to the topside atmosphere. Simple linear analysis suggests that a beam propagating through the magnetosphere will be stable to two-stream instabilities (Galvez and Borovsky, 1988), and a beam propagating into the ionosphere will be stable to resistive hose, ion hose and filamentation instabilities (Gilchrist et al., 2001). Simulations that track the beam from its source in the magnetosphere to its contact with the topside ionosphere are currently being carried out to quantify the effects of beam-plasma interactions as the beam moves through magnetic field and density gradients. Initial particle-in-cell simulation results,

supported by theoretical analysis, suggest no major effect of instabilities on the beam propagation (Kaganovich, private communication). Theory and simulation results will be reserved for a future publication.

AUTHOR CONTRIBUTIONS

ES assembled the INSPIRE team and, as Principal Investigator of the project, he coordinated research activities to ensure their relevance in addressing solutions to compelling magnetospheric problems, contributed to calculations of MHD magnetic field used to guide MeV electrons' trajectories, contributed to calculations of detectability thresholds of atmospheric effects measured with radar, and contributed to definition of beam properties for science applications. PP developed the loss cone calculations for MeV electron beams and developed the ballistic code for electrons. AP extended ballistic code for magnetic fields inferred from MHD global models. IK contributed with PP, AP, JJ, and DS developing ballistic codes and trajectory visualization tools. RM applied Monte Carlo techniques to measure atmospheric effects of MeV electron precipitation. JJ, MG-M, and KA developed algorithms to infer electron propagation properties in analytic magnetic field models. MN contributed to calculations of beam detection with radars and to definition of beam properties for science applications.

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The Beam Plasma Interactions Experiment: An Active Experiment Using Pulsed Electron Beams

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The 1970s and 1980s were heydays for using active electron beam experiments to probe some of the fundamental physical processes that occur throughout the heliosphere and in astrophysical contexts. Electron beam experiments were used to study spacecraft charging and spacecraft-plasma coupling; beam-plasma interaction physics; magnetic bounce and drift physics; auroral physics; wave generation; and military applications. While these experiments were enormously successful, they were also limited by the technologies that were available at that time. New advances in space instrumentation, data collection, and accelerator technologies enable a revolutionary new generation of active experiments using electron beams in space. In this paper we discuss such an experiment, the Beam Plasma Interactions Experiment (Beam PIE), a sounding rocket experiment designed to (a) advance high-electron mobility transistor-based radio frequency (RF) linear accelerator electron technology for space applications and (b) study the production of whistler and X-mode waves by modulated electron beams.

Keywords: active experiments, electron beam, wave-particle interactions, energetic particles, radiation belts, remediation, space physics

INTRODUCTION

Active space experiments using electron beams started in the 1970s primarily to study spacecraft charging effects (e.g., Mullen et al., 1986; Sasaki et al., 1986, 1988; Banks et al., 1990). In those experiments electron beams could produce controlled amounts of “artificial” charging in order to better understand the physical processes involved in spacecraft charging and neutralization and to investigate the effects of severe charging on spacecraft systems. Later electron beams were used to conduct a variety of innovative and successful active physics experiments involving beam-plasma interactions (e.g., Gendrin, 1974; Cambou et al., 1978, 1980), magnetic bounce and drift physics (e.g., Hendrickson et al., 1975, 1976; Winckler et al., 1975), and the generation of VLF wave emissions (e.g., Monson et al., 1976; Dechambre et al., 1980; Obayashi et al., 1982; Farrell et al., 1988; Neubert et al., 1988; Reeves et al., 1988, 1990a,b). Los Alamos National Laboratory also tested a neutralized H⁺ particle beam in the BEAR (Beam Experiment Aboard a Rocket) program as part of the US Strategic Defense Initiative.

Studies of wave-generation and wave-particle interactions using electron beams were of particular interest in the early days of active experiments. The experiments of the 1980's were able to demonstrate the ability to produce propagating electromagnetic waves; to identify that the strongest emissions were whistler-mode and; to establish a general agreement with analytic theory (Harker and Banks, 1985, 1987; Reeves et al., 1990b).

Theoretical work on beam-plasma-wave generation began in the 1960s and was further developed specifically for the active experiments program (e.g., Harker and Banks, 1983, 1985; Farrell et al., 1989). As with a physical antenna each beam pulse acts as a current source. The plasma through which the beam propagates responds according to the resonance condition:

$$\omega - k_{\parallel} v_{\parallel} = s \frac{\omega_{ce}}{\gamma} \quad (1)$$

where ω is the frequency of the mode, $k_{\parallel}$ is the wave vector along the background magnetic field (i.e., in the parallel direction), $v_{\parallel}$ is the parallel beam velocity, γ is the beam relativistic factor, ω_{ce} is the electron cyclotron frequency and s is an integer number with $s=0$ corresponding to the Landau resonance and $s \neq 0$ describing cyclotron resonance. The waves are emitted as Cherenkov radiation as described by e.g., Farrell et al. (1989). Harker and Banks (1985) calculated the whistler-mode wave power expected from a pulsed electron beam and Reeves et al. (1990b) found that the Spacelab 2 observations generally showed the predicted dependence on modulation frequency, duty cycle, and pitch angle. However, both the accelerator and receiver technologies available for the early electron beam experiments were quite limited. For example, the Spacelab 2 experiment could only operate at a single beam energy (1 keV) and current (100 mA). Only the modulation frequency and duty cycle could be varied. The wave receiver was a 1D analog audio recorder with a 10 kHz passband. No information on Poynting flux, wave normal angle, or polarization could be obtained.

Recently, Los Alamos National Laboratory and NASA's Goddard Spaceflight Center have been awarded a grant to conduct active experiments on beam-wave generation using state-of-the-art linear electron accelerators, wave receivers, and plasma instrumentation. The project is funded through NASA's Low Cost Access to Space (LCAS) sounding rocket program. The experiment is called the Beam Plasma Interactions Experiment, or Beam PIE, and is scheduled to launch in spring of 2021 nominally from Poker Flat, AK. In the remainder of this paper we discuss the experimental setup, objectives, and expected results based on theory, modeling, and simulation.

EXPERIMENTAL OBJECTIVES

Waves and wave-particle interactions play a critical role in some of the most important dynamics in space and astrophysical plasmas by mediating the exchange of energy between fields and particles. The Earth's radiation belts are a good example of such a system. In addition to the ULF wave-particle interactions that drive radial diffusion (and betatron/Fermi acceleration), plasmas injected from the magnetotail into the inner magnetosphere

form distributions that are energetically unstable. Depending on the nature of the plasma distributions and the ambient field and plasma conditions, those unstable distributions produce whistler-mode chorus, electromagnetic ion cyclotron (EMIC), magnetosonic, electron cyclotron harmonic, and other waves. Those waves, in turn, strongly affect the dynamics of the radiation belts. Whistler-mode chorus waves, for example, can strongly accelerate 100s keV "seed" electrons to MeV energies. These wave-particle interactions are considered to be the dominant radiation belt electron acceleration mechanism for at least a subset of events (Reeves et al., 2013; Thorne et al., 2013; Baker et al., 2014). EMIC waves can strongly pitch angle scatter radiation belt electrons and are candidates for rapid radiation belt losses (for at least some events and some energies) (e.g., Jordanova et al., 2008; Ukhorskiy and Sitnov, 2012; Usanova et al., 2014). Countless other examples can be found throughout the heliosphere: in the corona, solar wind, planetary magnetospheres, ionospheres, heliopause, and essentially every plasma system where we have wave and particle observations. Besides the scientific interest associated with the natural environment, wave-particle interaction physics can have very important practical applications such as radiation belt remediation (Carlsten et al., 2019).

There are many sources of free energy for wave generation and much study has been devoted to linear, quasi-linear, and non-linear instabilities that occur, naturally, in space [see e.g., Gary (1993)]. An alternative approach is a more active experimental technique, namely using an artificial electron beam to generate the waves. Accelerator-produced electron beams are "artificial" only in the sense that we can precisely control the characteristics of the beam to produce waves with equally precise and testable characteristics.

The objective of the Beam Plasma Interactions Experiment is to discover and characterize fundamental wave-particle interactions by generating waves using a modulated energetic electron beam, characterizing the wave properties to test theoretical and model predictions. As a secondary objective we will determine if the beam-generated wave fields are strong enough to produce measurable scattering of ambient ionospheric electron populations. The main specific objectives of Beam PIE are to:

- Demonstrate, for the first time, advanced RF linear electron accelerator instrumentation for space experiments.
- Quantitatively test theories of how energetic electron beams couple to plasmas to stimulate whistler-mode radiation.
- Discover and characterize how energetic electron beams couple to plasmas to stimulate propagating R-X-mode radiation.

BEAM PIE—THE BEAM PLASMA INTERACTIONS EXPERIMENT

Experimental Concept of Operations

Beam PIE will utilize a "mother-daughter" rocket configuration in which one rocket segment, the "accelerator" will house the electron beam and power systems and the other section, the

“receiver” will house the fields, waves, and particle detectors. The payload layout in launch configuration and placement of the instruments on the two payloads are shown schematically in **Figure 1**.

After achieving operational altitude and immediately after engine cutoff an attitude control system (ACS) will be used to orient the payloads such that they will be aligned $\pm 2^\circ$ to the magnetic field at apogee which achieves $\pm 5^\circ$ for all altitudes above 300 km. The orientation of the magnetic field is known much more precisely from models and previous rocket missions but some margin of error is built into the mission success criteria. At all times through flight, magnetometers on both the accelerator and receiver payloads will provide more precise knowledge of the orientation with respect to B and therefore beam pitch angle and receiver “spin tone.”

The accelerator and receiver segments will be spring-separated to place the accelerator segment on a higher altitude trajectory (apogee ~ 500 km) and the receiver segment on a lower-altitude trajectory with ~ 1 km peak separation as achievable by the spring system (**Figure 2**). With the rocket body oriented along B the separation creates a V primarily in the B direction. Thus, separation of the accelerator and receiver payloads is roughly field-aligned with minimal separation in the perpendicular to B direction. GPS receivers on both payloads will provide knowledge of the payload separations both along and perpendicular to B.

The electron beam on the accelerator segment will be aligned with the rocket body and therefore also directed downward along the magnetic field line. The beam will be operated in a sequence of beam “firings.” The Beam PIE linear accelerator is capable of producing beams with energies ranging from ~ 10 –50 keV. The accelerator electronics are capable of modulating the beam at frequencies from a few Hz up to 1 MHz. Modulation frequencies of ~ 2 –25 kHz optimize the generation of whistler-mode waves. Coherence effects maximize power at the beam modulation frequency and harmonics thereof (Harker and Banks, 1985). In the R-X mode, waves are generated at frequencies between the plasma and upper hybrid frequency

regardless of the beam modulation frequency but coherence effects favor very short beam pulses (see **Figures 10, 11**). Using frequencies of 1 MHz and duty cycles $< 10\%$ produces pulses of < 100 ns which should concentrate nearly all the wave power in the R-X mode. Theory and modeling predict that both the division of power between the whistler and R-X-mode waves, and the characteristics (power, frequency, wave vector, polarization, etc.) of each wave mode, should be a strong function of the beam parameters and thus a sensitive test of our understanding of beam-plasma-wave interactions (Expected dependencies are discussed further in section Expected Results from Theory, Modeling, and Simulation).

The receiver payload includes a 3-axis electric field measurement, 3-axis search coil, fluxgate magnetometer, and a full waveform capture digital receiver that will exquisitely characterize the waves generated by the beam from the accelerator payload. The receiver payload will also measure the parameters needed to calculate the wave dispersion relation: the background magnetic field, electron temperature, and absolute electron density.

It is not our objective to measure the electron beam itself on the receiver payload. Although the plasma instrument is capable of that measurement, the beam is guided by the magnetic field and the gyroradius of the beam is extremely small relative to the size of the receiver payload. Therefore, only very precise (and fortuitous) magnetic conjunction would allow direct detection of the beam. Rather, the experiment is designed to measure beam-generated waves and scattering of ambient electrons—neither of which requires precise magnetic conjugacy.

The phase velocities of both whistler and R-X-mode waves that are generated by the electron beam are equal to the beam velocity. The group velocity of the whistler waves is in the same direction as the beam itself (i.e., downward). Those waves are fast and propagate mostly field aligned. On the other hand, the R-X-mode waves have a much smaller group velocity, oriented opposite to the beam over a wide range of angles that strongly depends on frequency. Thus, in order for the instruments in the

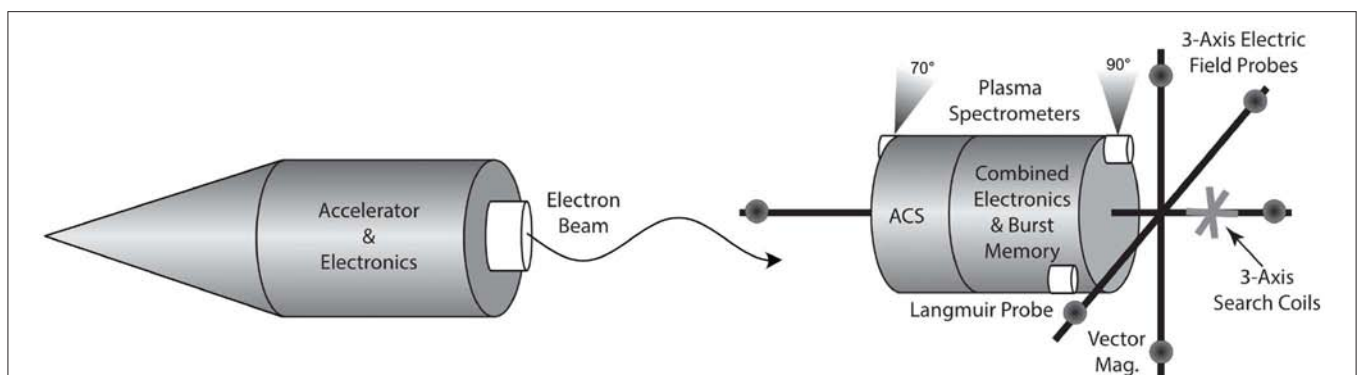


FIGURE 1 | Schematic of the Beam PIE payload and instrument locations. The Accelerator Payload contains only the accelerator and associated electronics. After payload separation it will be aligned roughly along the magnetic field and boosted to a higher altitude. The Receiver Payload consists of 3-axis electric and magnetic field probes, a wave digitizer/receiver, a plasma spectrometer, a Langmuir probe, a vector magnetometer and associated electronics [see sections DC and Wave Electric Field Detectors–Energetic Electron Spectrometer (APES)]. An attitude control system (ACS) will be used to achieve a spin rate for the receiver payload of ~ 1 Hz. The same ACS will orient the payloads such that they will be aligned $\pm 2^\circ$ to the magnetic field at apogee which achieves $\pm 5^\circ$ for all altitudes above 300 km.

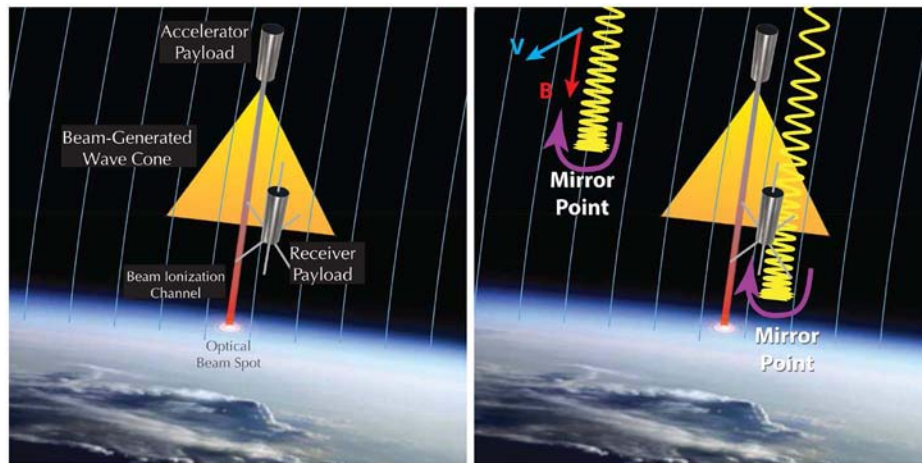


FIGURE 2 | Schematic of the Beam PIE concept of operations. The Accelerator Payload directs the electron beam down the magnetic field line. The beam-generated waves spread out in a cone of wave power. The waves are detected and characterized by the Receiver Payload flying at a somewhat lower altitude. In the region where waves are present, they resonantly pitch angle scatter ambient ionospheric electrons which lowers their magnetic mirror points and increases the flux of upward going electrons detected at the Receiver Payload.

receiver payload to detect the R-X mode waves generated by the beam injected from the accelerator payload, the receiver must be in or above the wave generation region. As is discussed above, the two payloads are spring-separated so that their distance increases with time, reaching a peak separation of about 1 km during the ~ 400 s duration of the experiment. The minimum separation during operation is on the order of ~ 100 m, since the beam firing sequence starts tens of seconds after the two payloads start to separate. The power of the waves emitted by the beam is maximized when the characteristic size of the beam pulse along the field is less or comparable to the wavelength of the mode to be excited. The expansion of the beam pulses imposes a constraint on the size of the emission region (see Section Expected Results from Theory, Modeling, and Simulation), which is stronger for R-X modes due to their smaller wavelength. Preliminary calculations indicate that the beam may continue to efficiently excite R-X mode waves over distance of several hundred meters (Delzanno and Roytershteyn, 2019), implying that the receiver payload will travel through the R-X wave generation region (which is created almost instantaneously since the beam pulses travel at a fraction of the speed of light) as its distance relative to the main payload increases. Furthermore, R-X mode waves created below the receiver payload will move upward and at least some of them will be captured by the receiver, possibly in the far-field depending on the actual cross-field separation between the two payloads. Since whistlers have significantly larger wavelength, the size of the whistler-wave-generation region is significantly larger and the receiver will be comfortably inside it. Further, the receiver could also capture whistler waves generated above the receiver payload and moving downward. We note that, in general, the area of the beam itself is only a few gyro-radii but the area over which waves can be detected is much larger. For waves propagating $\pm 20^\circ$ with respect to B, a 1 km accelerator-receiver separation produces cone of

wave power ~ 700 m in diameter. Larger propagation angles, of course, produce a larger area of radiated power. Therefore, while the accelerator and receiver should be roughly field-aligned, considerable cross-field separations are not at all problematic and can in fact allow measurements in the far-field.

It is well-known that whistler mode waves pitch angle scatter electrons through Langmuir and gyro-resonant wave-particle interactions but the electron scattering by R-X-mode waves has not been tested in space. If the R-X-mode waves are sufficiently strong they would be expected to very efficiently pitch angle scatter ambient (background) ionospheric electrons and we will also check for this effect. Because of their polarization R-X-mode waves pitch angle scatter electrons that are traveling in the same direction as the wave phase velocity— in this case that means electrons moving down the field lines. The accelerator and receiver payloads will both operate at ~ 300 – 500 km, well above the nominal atmospheric absorption altitude of ~ 100 km for 10 s keV electrons. Therefore, the pitch angle distribution will have a large atmospheric loss cone and very strong anisotropies when comparing the upward and downward directed hemispheres. When the beam-generated waves are off, few particles will mirror below the receiver payload and therefore few particles should be measured moving up the field. When the beam is on, we will look for pitch angle scattering of ambient ionospheric electrons by looking at a change in the flux of upward-going electrons. A change in upward-going electrons could indicate that the waves are scattering electrons to lower mirror altitudes or that these electrons could have been precipitated (**Figure 2B**). To make this measurement we include two plasma spectrometers capable of measuring 0–30 keV electrons [see section Energetic Electron Spectrometer (APES)]. The receiver payload will be spinning and oriented with the spin axis aligned to the Earth's magnetic field so one plasma spectrometer will be oriented at 90° to measure locally mirroring electrons and one will be

mounted facing “downward” to measure upward going electrons that have mirrored below the receiver payload. The “downward-looking” spectrometer will be oriented 70° to the spin axis which will measure electrons mirroring at ~ 100 km altitude (Whistler mode waves resonate with electrons propagating in the opposite direction as the waves and also resonate with much higher energy electrons so whistler mode scattering is not readily measured with this experimental set up).

Beam PIE Instrumentation

As discussed above, Beam PIE is a standard “mother-daughter” rocket configuration consisting of a main payload and a subpayload. The main “receiver” payload will house the fields, waves, and particle detectors and will be equipped with an ACS (attitude control system). The “daughter,” or “accelerator” subpayload, located forward of the main payload, will house the electron beam accelerator and power system (Figure 1).

Linear Electron Accelerator

The electron accelerator is shown schematically in Figure 3. The electron beam is generated in a DC ~ 10 -keV “electron gun” and injected into a single 5-GHz RF cavity which can accelerate the beam an additional 40 keV for a total nominal energy range of ~ 10 –50 keV.

When operating at 10 keV, the electron gun will produce 20 mA of current. A bare accelerator system produces a beam with $\sim 100\%$ E/E which is too large for efficient wave generation. Therefore, we will use a chicane magnet at the beam aperture to reduce the beam E/E to $\sim 10\%$. This also reduces the emitted beam current to $\sim <2$ mA which is sufficiently small that no significant spacecraft charging is produced.

The 10 keV electron gun is commercial off the shelf (COTS). The accelerator cavity uses a novel LANL design that adapts common laboratory RF linear accelerator (linac) components to be suitable to space applications. One novel feature of the design is the use of high-electron mobility transistors (HEMTs) to energize the RF cavity (Figure 4). HEMTs greatly reduce power consumption and associated waste heat.

A notional electron beam pulse format is also shown in Figure 3. Nominal operations consist of a sequence of $\frac{1}{2}$ second beam pulses separated by $\frac{1}{2}$ second when the beam is off and no waves are being generated. Each $\frac{1}{2}$ s of beam on time can utilize a different combination of beam energy, RF frequency, and beam duty cycle. Additionally, the range of rocket altitudes provides a range of background plasma and magnetic field conditions allowing a wide range of beam-plasma-wave interaction conditions to be investigated. The $\frac{1}{2}$ s intervals when the beam (and waves) are off allows unambiguous separation of beam-generated waves from naturally occurring wave conditions.

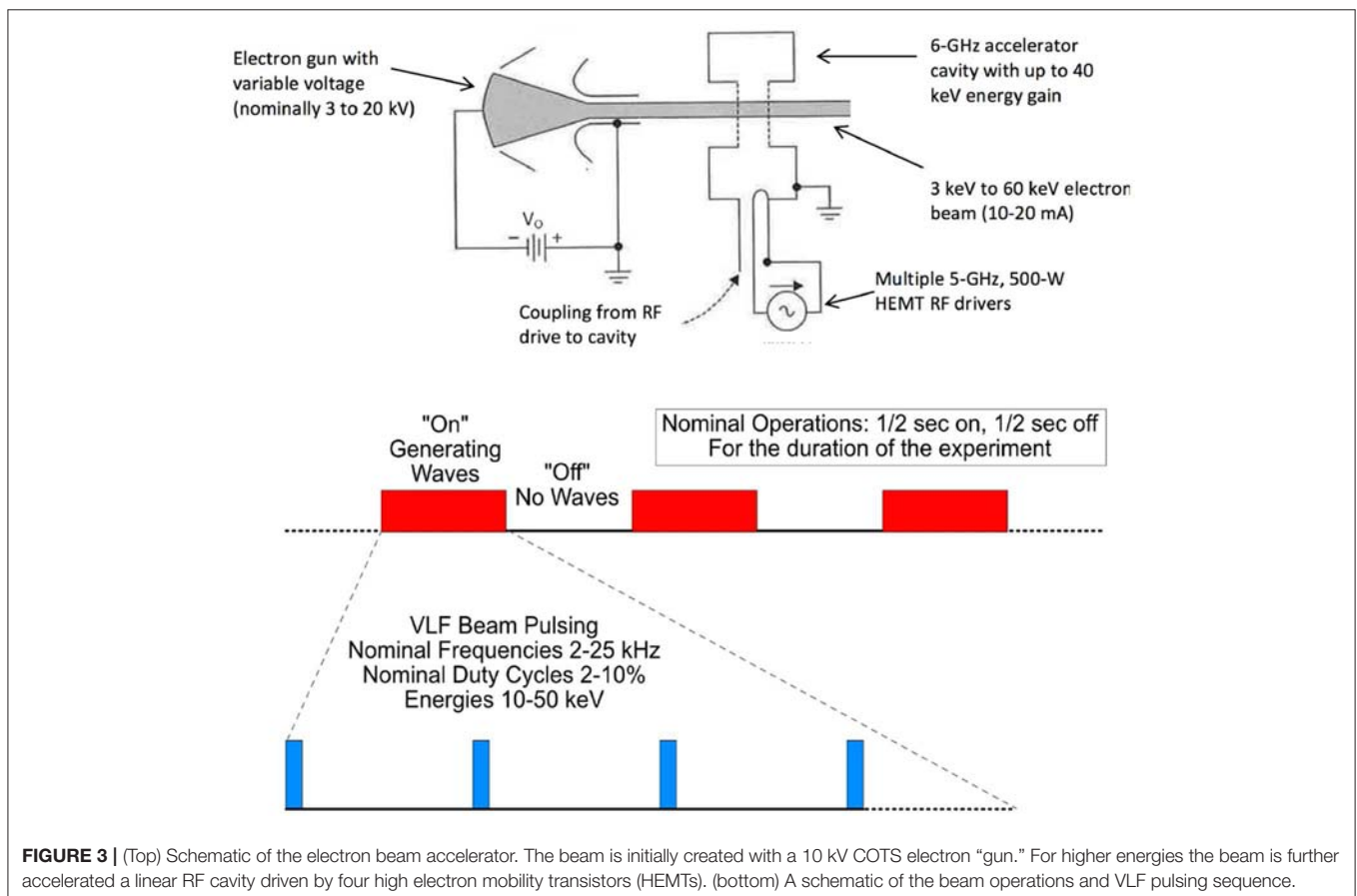
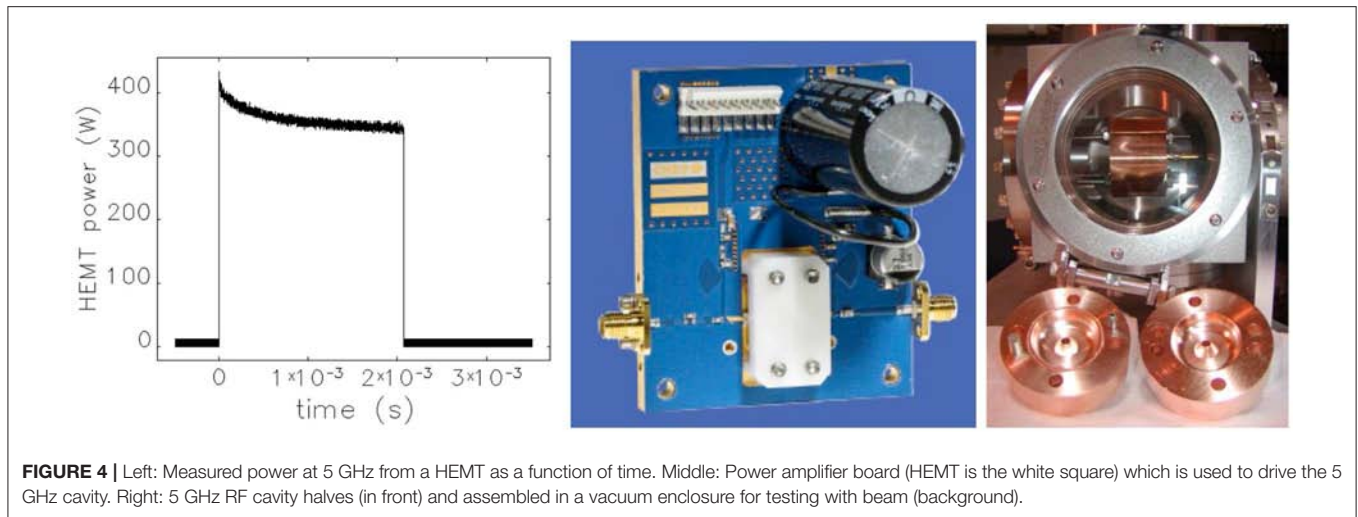


FIGURE 3 | (Top) Schematic of the electron beam accelerator. The beam is initially created with a 10 kV COTS electron “gun.” For higher energies the beam is further accelerated a linear RF cavity driven by four high electron mobility transistors (HEMTs). (bottom) A schematic of the beam operations and VLF pulsing sequence.



DC and Wave Electric Field Detectors

The Beam PIE wave receiver will gather measurements of both ambient waves and plasma structures as well as plasma waves excited by the electron beams. The DC and AC vector electric field will be measured using the standard double probe technique (e.g., Pfaff, 1996). In this manner, spherical sensors with embedded pre-amps will be extended on three independent, tri-axial double probes. This configuration provides a full, three dimensional electric field measurement that will completely parameterize the vector electric field (DC to 5 MHz) including DC and wave electric fields parallel to the magnetic field direction (Figure 5).

The double probes include inner spheres (situated 0.5 m inboard of the outer spheres) to serve as multiple baseline electric field detectors or spaced receivers (Pfaff and Marionni, 1998). These “double-double” probes are similar to ones flown successfully in the auroral E-region Rocket/Radar Instability Study (ERRIS) (rockets 21.097 and 21.100) and Cusp Transient Features Campaign (rockets 36.152, 36.153). Measurable phase shifts from these separated receivers not only establish the electrostatic nature of any ELF/VLF wave modes associated with the electron beams, but also provide a measure of their wavelength and phase velocity.

The wave receiver also includes an HF channel to observe the presence of any waves near the electron plasma frequency, such as HF Langmuir waves. The electronics will return continuous FFT power spectra of the ambient plasma environment. Furthermore, a burst memory will gather vector AC fields (three components) sampled at 10 Msample/sec each synchronized with the electron beams with ample time prior and after the actual beam discharge. Importantly, the burst waveform capture capability and dedicated telemetry system allow the opportunity for discovery of phenomena outside of the primary experimental objectives including instabilities and potential non-linear effects.

Langmuir Probe

A fixed-biased Langmuir probe will be flown in order to observe the electron plasma number density and its fluctuations. The



FIGURE 5 | Electric field probes as deployed in sounding rocket 21.116.

Langmuir probe will be oriented perpendicular to the spin axis and magnetic field direction to minimize spin effects. In addition to pre-launch calibration curves and theory, the Langmuir probe will be normalized using simultaneous ground-based Poker Flat

Rocket Range ionosonde data, as well as plasma wave data where applicable.

Vector Magnetometer

The Goddard Space Flight Center will furnish a vector fluxgate magnetometer similar to that flown on numerous sounding rocket experiments to measure the currents associated with the aurora, the Sq current system, and the equatorial electrojet. This magnetometer is a commercially procured Bartington type (or equivalent) for which 18-bit A/D converters will be built at the Goddard Space Flight Center. The electronics are part of the electric field electronics. These instruments provide tri-axial measurements to a resolution better than ± 1 nT. This performance is sufficient to detect VLF whistler mode waves as well as the potential effects of field-aligned currents. The magnetic field component of MHz R-X mode waves will not be measured in this experiment.

Search Coil

The Search Coil unit is an AC magnetic field sensor capable of detecting the B-field vector of an electromagnetic plasma or radio wave. Each coil has many thousands of turns to convert a wave's dB/dt into an output sensor voltage between 10 Hz and 100 kHz. The analog output $V(t)$ from each sensor can then drive an ADC for inclusion in the data stream for return and subsequent spectral analysis. **Figure 6** shows a set of flight units built by Goddard Space Flight Center for the Air Force's DSX mission (Scherbarth et al., 2009).

Combined Electronics and Burst Memory

The electronics for the combined DC/Wave electric field instrument and the fluxgate and search coil magnetic field instruments, as well as the Langmuir probe, will be integrated into a combined electronics assembly built by Goddard. Similar unites have been flown on numerous previous experiments. The Goddard group will also provide a burst memory for the payload that will be synchronized with the electron beams, as was successfully carried out as part of the fields experiment for the APEX high density beam releases (see Pfaff et al., 2004).

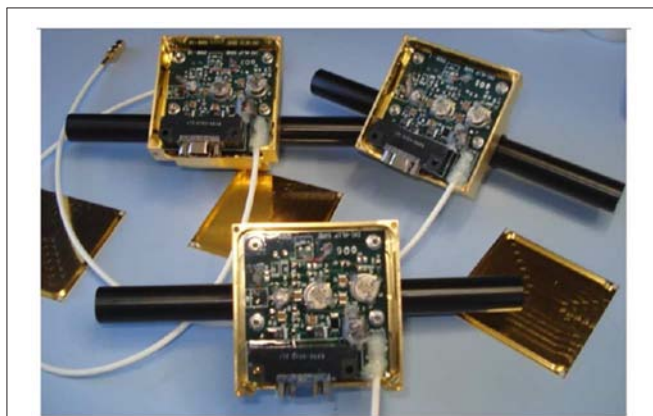


FIGURE 6 | Individual search coil units with per-amp covers removed.

The burst memory will gather precursor data prior to the beam activation, to ensure that the fastest particles and wave modes associated with the release itself will be captured. The burst memory will record HF and MF vector data gathered by the electric field detectors.

Energetic Electron Spectrometer (APES)

The preceding instruments, as described, are all required for the primary Beam PIE science objectives—generation of whistler and R-X-mode waves with a novel linear electron accelerator. We also note that the previously-discussed receiver payload instruments can operate in any orientation with respect to the magnetic field and can operate equally well on a non-spinning platform.

The secondary objective of Beam PIE, however, is to study the effect of the waves on ambient electrons. Specifically we investigate pitch angle scattering by R-X mode waves.

The Beam PIE energetic electron spectrometer (known as the Acute Precipitating Electron Spectrometer, APES) uses magnetic deflection to measure the locally-mirroring and upward-going electrons with high cadence over a 150 eV to >30 keV energy range. The APES field of view is $10^\circ \times 10^\circ$. One APES spectrometer is oriented at $\sim 90^\circ$ to the ambient magnetic field to measure locally-mirroring electron. The other is oriented at 70° to the spin axis in order to measure upward-going electrons that mirror below the rocket at ~ 100 km altitude. An increase in the ratio of upward going electrons to locally-mirroring electrons indicates pitch angle scattering by the R-X mode waves.

APES [described in detail in Michell et al. (2016)] uses a micro-channel plate (MCP) detection system with 50 discrete anodes (energies). Ray tracing analysis of the magnetic deflection system to be used is presented in **Figure 7** next to a photograph of the APES instrument that flew on the GREECE mission.

Ground-Based Diagnostics

To the extent feasible, radar and optical ground-based instrumentation will be used to measure the properties of the ambient plasma before and during the beam injection and to remotely measure the ionospheric effects of the primary beam energy deposition and enhanced precipitation from wave-particle interactions. The effects that we will be looking for is the electron density (N_e) perturbation over background, measured with incoherent scatter radar and optical emissions using ground-based imagers.

The Poker Flat Incoherent Scatter Radar (PFISR) is operated by SRI International on behalf of the National Science Foundation and it will be operated during the rocket flight. PFISR is a modular, UHF phased-array capable of beam steering on a pulse-to-pulse basis (Nicolls et al., 2007). PFISR operations will start at least 2 h before and continue at least 2 h after the rocket's launch window.

The radar's mode of operation for the Beam PIE rocket flight will consist of multiple beams with orientations selected to measure the ambient N_e and convection along the foot-point of the rocket's trajectory and its vicinity, as well as the perturbation N_e caused by the energy deposition of the electron

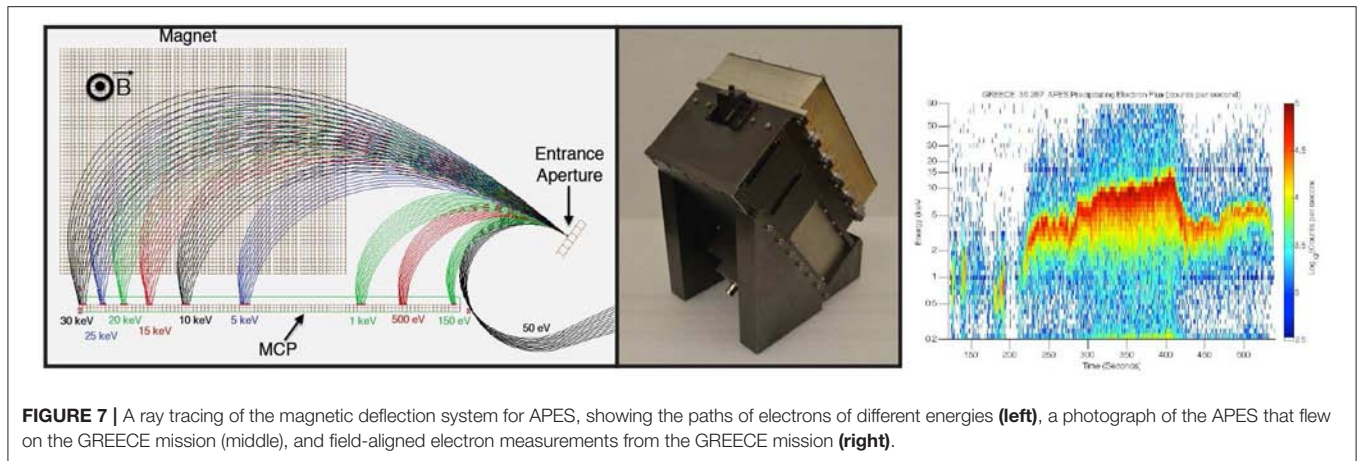


FIGURE 7 | A ray tracing of the magnetic deflection system for APES, showing the paths of electrons of different energies (left), a photograph of the APES that flew on the GREECE mission (middle), and field-aligned electron measurements from the GREECE mission (right).

beam. Measurements of the electron density perturbations as a function of beam energy and current will provide diagnostics of beam propagation from the accelerator payload and energy deposition into the atmosphere. Detection of optical emissions from the interaction of the beam with the atmosphere provide opportunities for additional beam diagnostics. The optical emissions are similar to naturally-occurring aurora but, with this active experiment we have precise knowledge of beam energies, currents, and frequencies [e.g., Marshall et al. (2014)].

EXPECTED RESULTS FROM THEORY, MODELING, AND SIMULATION

Expected Results—Wave Generation

The primary objectives of the Beam Plasma Interactions Experiment are to test our understanding of wave generation through beam-plasma interactions. We will independently test the generation of whistler-mode and R-X-mode waves as well as the partitioning of energy between the modes. Wave diagnostics include wave spectral power density, polarization, ellipticity, and Poynting flux. While the space physics community has come to understand the power of such measurements, what makes Beam PIE unique is the ability to do active, controlled experiments with unprecedented flexibility in pulsed electron beam parameters. We have the ability to independently vary beam energy, pulse frequency, and duty cycle. In addition, the trajectory of the rocket naturally samples different background plasma conditions (density, temperature and magnetic field strength) that determine the plasma wave dispersion relation.

For wave generation, the starting point is the theory of Harker and Banks (Harker and Banks, 1983 labeled “HB” in the following; 1985; 1987) which solves the cold plasma dispersion relation to obtain the wave amplitudes for a sequence of beam pulses with finite length (labeled “ l_p ”) along the direction of motion, while the pulses are infinitesimal in the perpendicular direction. These pulses are assumed to move with a constant velocity and pitch angle α relative to the background magnetic field and the response of the plasma to the beam pulses is calculated in the framework of cold-plasma linear theory.

The HB theory was developed in the 1980’s, stemming from earlier work that calculated the radiated power from a point charge (McKenzie, 1963; Mansfield, 1967) and was used for the interpretation of ionospheric electron-beam experiments such as Spacelab 1 and 2 (Gurnett et al., 1986; Bush et al., 1987; Farrell et al., 1988, 1989; Reeves et al., 1988, 1990a,b; Farrell and Goertz, 1990). It is used here as a reference and, later in this section, we discuss its limitations and the physical effects that will need to be incorporated for more accurate predictions of waves generated by Beam PIE.

In order to obtain realistic parameters for Beam PIE, we have used the International Reference Ionosphere (IRI) model for the month of March 2009 at 0 Local Time and for altitudes between 300 and 500 km. The corresponding average density is $n = 3.8 \cdot 10^4 \text{ cm}^{-3}$. The average magnetic field for the same altitudes at the Poker Flat (Alaska) launch site is $B_0 = 4.7 \cdot 10^{-5} \text{ T}$. These parameters give the ratio of the electron plasma frequency (ω_{pe}) to the cyclotron frequency (ω_{ce}) equal to $\omega_{pe}/\omega_{ce} = 1.33$. In what follows, we assume hydrogen ions.

Generation of Whistler Waves

Figure 8 shows the power spectral density obtained from HB in the whistler regime ($\omega \leq \omega_{ce}$), for three beam energies equal to 14, 34, and 54 keV. The former value corresponds to the maximum energy obtained from the DC electron gun, while the latter value is the maximum energy achievable after the RF accelerator cavity. The calculation is performed for a single beam pulse of length corresponding to the pulse period $t_p = 100$

ns and the power spectral density is in units of $[\frac{q_p^2 \omega_{ce}}{\epsilon_0 c}]$ (where q_p is the pulse charge, ϵ_0 is the permittivity of vacuum and c is the speed of light). We have also used a beam pitch angle $\alpha = 1^\circ$ to account for inaccuracies in beam pointing relative to our target of injection aligned to the background magnetic field and only computed contributions due to the Landau resonance (as appropriate for a field-aligned beam). Note also that a finite pitch angle is necessary to obtain a finite total radiated power (McKenzie, 1963). Figure 8 (left) shows that the whistler signal breaks into two distinct frequency bands for $E = 34 \text{ keV}$ and shrinks considerably in frequency for $E = 54 \text{ keV}$. Consistently,

the whistler radiated power drops by two orders of magnitude for energies above $E \sim 35$ keV.

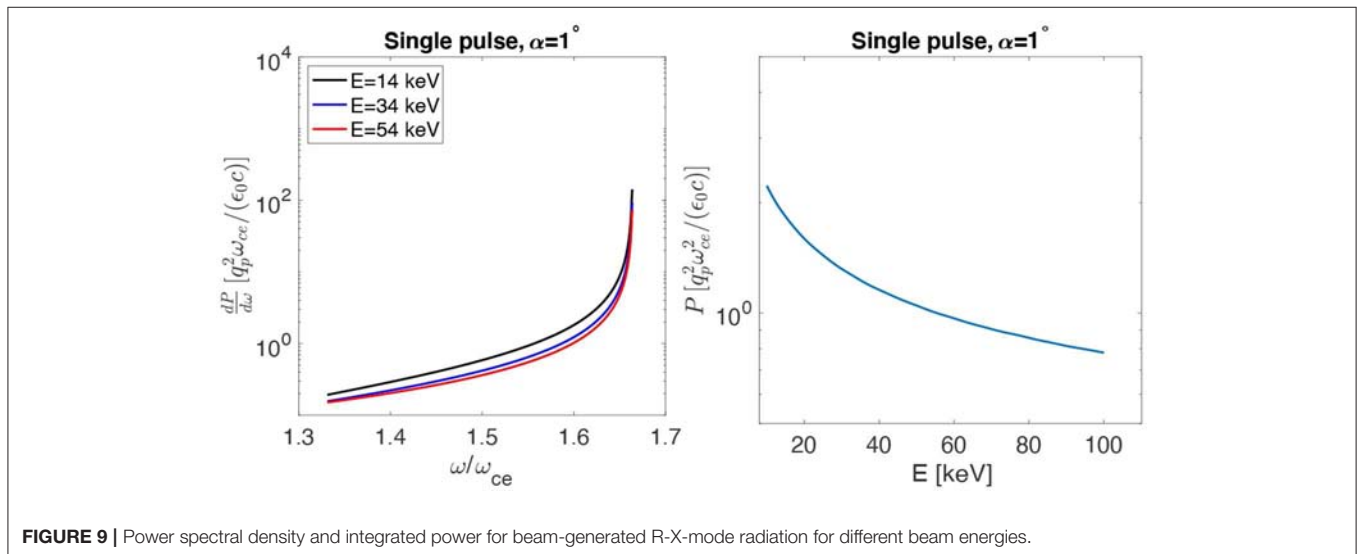
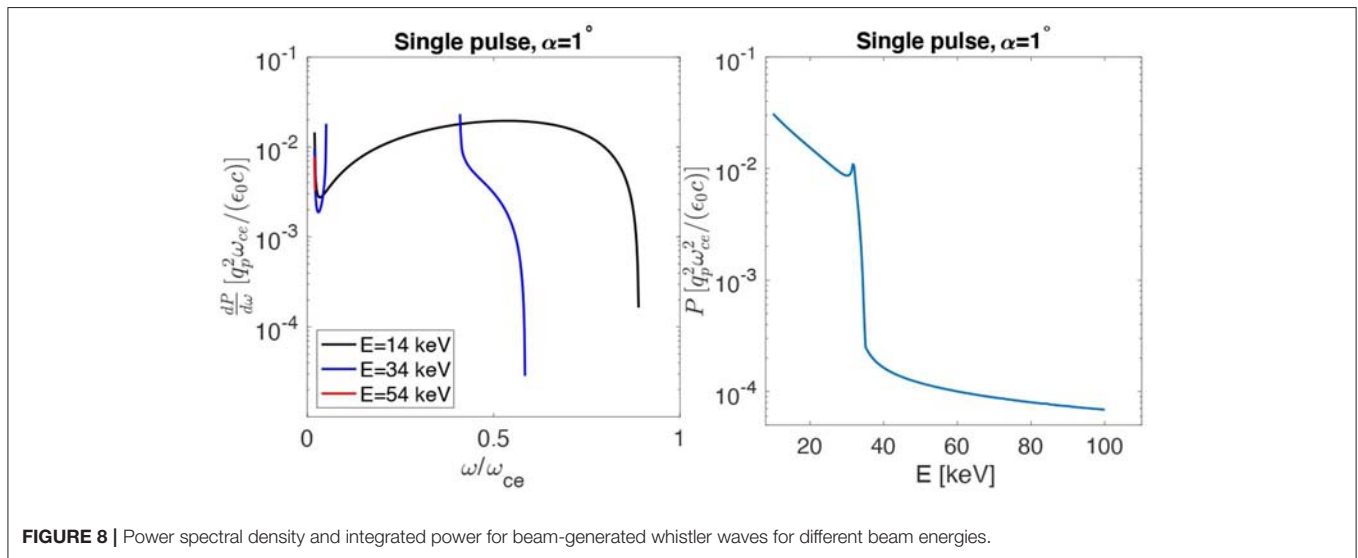
Generation of R-X-Mode Waves

The whistler branch is not the only regime where the beam can couple with a magnetized plasma. **Figure 9** (left) shows the power spectral density vs. frequency with the same format of **Figure 8**, extending the frequency range from $\omega = \omega_{pe}$ to the upper hybrid frequency $\omega_{uh} = 1.64\omega_{ce}$. Comparing against **Figure 8** shows that the R-X-mode wave signal is quite similar for the beam energies considered and, most important, it can be orders of magnitude stronger than the whistler mode signal. The total radiated power in **Figure 9** (right) shows a decreasing trend with beam energy but is several orders of magnitudes higher than the power radiated in the whistler range for all cases considered. These results indicate that highest radiation and, hence, highest beam-plasma coupling may be achieved through the R-X-mode.

Beam Operations for Wave Generation

While we have demonstrated that both whistler and R-X-modes can be excited by an electron beam, a primary objective of Beam PIE is to use controlled experiments to quantitatively test our understanding of beam-plasma-wave interactions. The beam parameters that can be user-selected are beam energy (i.e., velocity), modulation frequency, and pulse duty cycle. We have designed nominal beam operation modes to test each of these parameters.

An example of the predicted energy partitioning between whistler and R-X-mode waves is shown in **Figures 10, 11**, together with more detailed examples of wave dependence on beam parameters. **Figure 10** is obtained for beam energy $E = 14$ keV, pitch angle $\alpha = 1^\circ$, 100 pulses of length corresponding to 100 ns and varying the duty-cycle between 5 and 20%. This corresponds to a beam modulation between 500 kHz and 2 MHz, which targets the R-X mode. Indeed, despite the fact that both

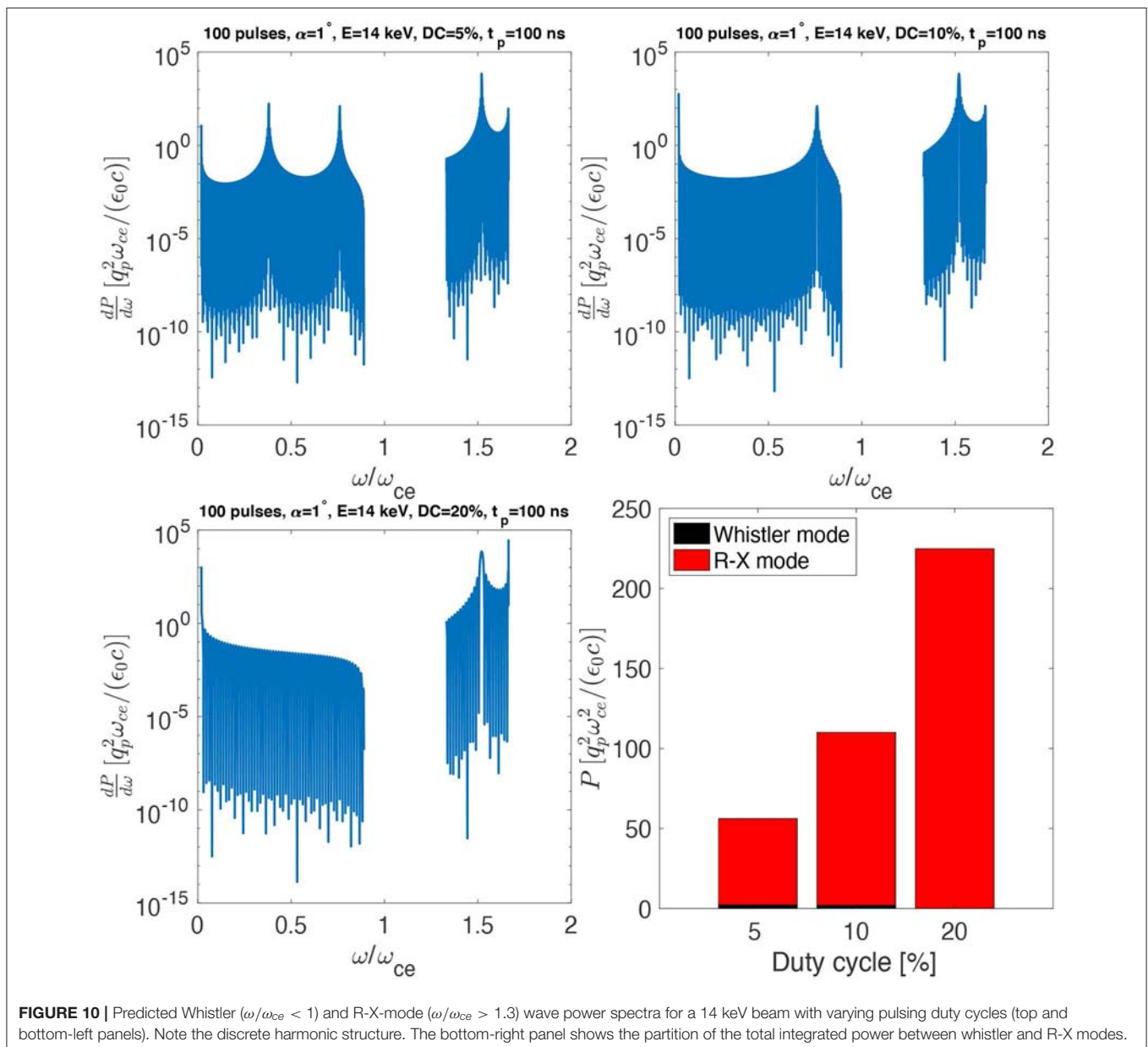


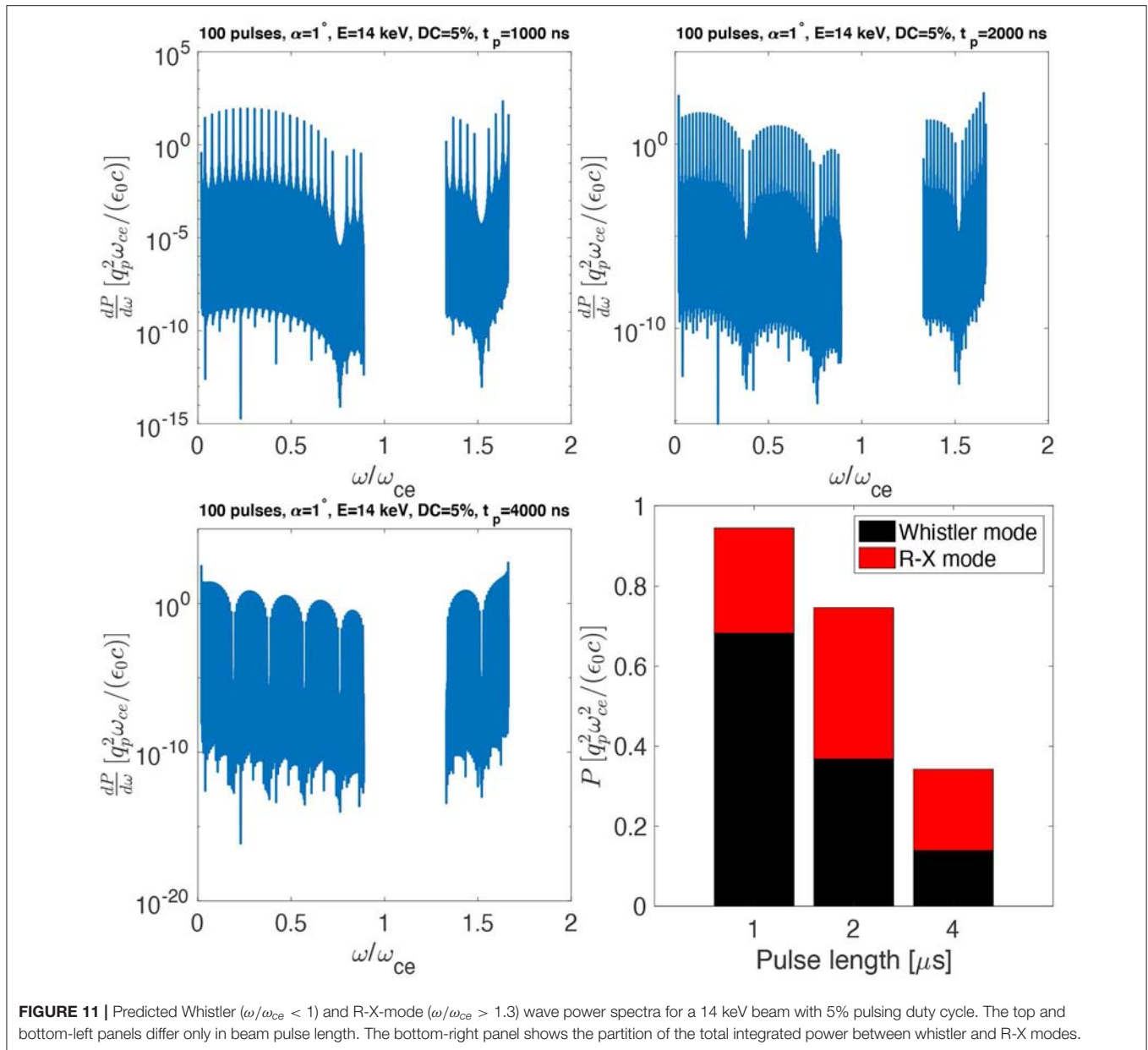
whistler and R-X modes are generated, the total radiated power is overwhelmingly in the R-X modes. **Figure 11** is obtained for the same parameters of **Figure 10**, except that the duty-cycle is fixed at 5% and the pulse length is varied between 1 μ s and 4 μ s. Since the duty-cycle is constant, varying the pulse length implies a modulation between 12.5 kHz and 50 kHz, i.e., in the whistler regime. Unlike the case in **Figure 10**, in this case the radiated power in whistler and R-X mode waves is comparable. Note that the total normalized radiated power in **Figure 11** is lower than that in **Figure 10** because it is plotted in units of $[q_p^2 \omega_{ce} / (\epsilon_0 c)]$. For the same beam current, a 1 μ s beam pulse has 10 times larger pulse charge than a 100 ns beam pulse, implying that in reality the total radiated power in dimensional units is comparable in the two cases. In each of these cases both whistler and R-X-mode

waves are generated. We note, however, that the partitioning of wave energy, the spectral shape in each mode, and the harmonic structures are all strongly frequency dependent providing a very sensitive test of the theory of beam-plasma-wave interactions.

Simulation Support for Beam PIE

The HB theory provides a good starting point to estimate the response of a magnetized plasma to a pulsed electron beam. However, it has a number of limitations that need to be addressed to improve the accuracy of the Beam PIE predictions. In particular, the theory assumes an infinitesimal transverse direction of the beam and is based on cold-plasma theory. These approximations imply that there are resonances at the lower- and upper-hybrid frequencies where the power spectral density





diverges, even though the total radiated power remains finite (for non-zero pitch angles). The behavior of the power spectral density near resonances affects the total radiated power and the partition between whistler and R-X modes—particularly the power in the R-X modes since the power spectral density is monotonically increasing toward the upper-hybrid frequency, c.f. **Figure 9**. There are several effects that can regularize the power spectral density (finite transverse beam size, thermal effects, non-linear effects and collisions) around resonances and that need to be properly considered. Furthermore, the HB theory does not take into account beam dynamics nor the potential feedback between the plasma and the beam. The beam dynamics are important because the beam pulses can radiate efficiently (i.e., coherently) in certain wavelengths only if the longitudinal

extension of the beam pulse is smaller/comparable to the excited wavelengths and this affects the longitudinal extension of the radiated wave field. The beam dynamics are in itself fairly complex since the beam can expand longitudinally due to its space-charge, while oscillating transversely due to the combined effects of space charge and the Lorentz force. How much the beam pulse charge is neutralized by the background plasma also affects the beam dynamics. Last, the feedback between beam and plasma is a possible source of instability, whose effect on the radiation pattern needs to be evaluated. Generally speaking, earlier work showed that the small, finite-size beam radius decreases the growth rate of the electrostatic two-stream instability (Galvez and Borovsky, 1988) and early active experiments showed that beams could propagate long-distances in space (Winckler, 1982). A

Spacelab experiment, on the other hand, showed radiation levels consistent with coherent Cherenkov emission and attributed it to beam bunching due to the two-stream instability (Farrell et al., 1989). Note that the criteria for beam instability is expected to be $k_{\parallel} V_b \sim \omega_{pe}$, where V_b is the beam velocity (Gary, 1993). Since this is approximately the same criterion for R-X mode emission, it gives a pulse length for instability comparable to that where coherent effects will quench radiation into R-X modes. Thus, using short (100 ns) pulses maximizes radiation into R-X modes but also prevents the development of the instability. Longer pulses (either directly from the accelerator or elongated by pulse expansion via space-charge effects) could potentially become unstable and create bunches and this regime will also be tested by Beam PIE.

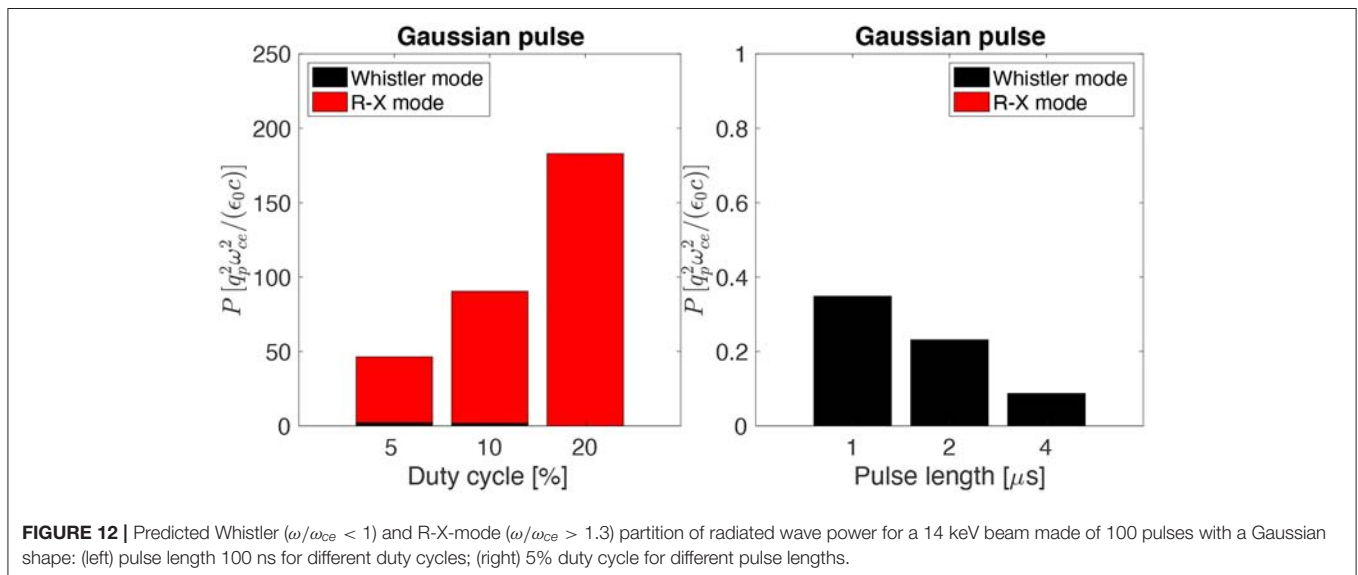
As a first step to address the limitations just discussed, we have modified the HB theory to account for pulses of finite transverse size, assuming pulses with a Gaussian shape characterized by a width along and perpendicular to the magnetic field, $l_{\parallel}$ and $l_{\perp}$, respectively. While a cylindrical shape is characteristic of pulses when they leave the RF accelerator (B. Carlsten, personal communication), a Gaussian shape might be more appropriate on longer time-scales of beam dynamics. We have applied the finite-transverse-size HB theory to the same case presented in **Figure 10**, for a beam with $l_{\perp} = 0.1$ m and $l_{\parallel} = 3.4$ m. The results for the partition of the radiated power between whistler and R-X modes vs. duty cycle are shown in **Figure 12** (left), where one can see that the radiated power is still dominated by R-X modes but it is lower by ~ 20 – 30% than what shown in **Figure 10**. **Figure 12** (right) shows the applications of the finite-transverse-size HB theory to the same case presented in **Figure 11**, for a beam with $l_{\perp} = 0.1$ m and $l_{\parallel} = 34$ m (pulse length = $1 \mu\text{s}$), $l_{\parallel} = 69$ m ($2 \mu\text{s}$) and $l_{\parallel} = 138$ m ($4 \mu\text{s}$), for a 5% duty cycle. While there is small reduction in the whistler radiated power, one can see that radiation in R-X modes is completely quenched (c.f. **Figure 11**) since the pulse lengths are now larger than the characteristic wavelength of the modes. These

results emphasize the importance of treating beam dynamics and the pulse evolution in the determination of the beam-generated wave field.

Furthermore, in order to address the issue of beam dynamics and stability, we have adopted a two-pronged approach. First, we are performing simulations of wave emission by a pulsed beam of a given shape with a highly accurate three-dimensional Vlasov code called the Spectral Plasma Solver (SPS) (Delzanno, 2015; Vencels et al., 2016; Roytershteyn and Delzanno, 2018). Second, we are performing simulations of beam pulse dynamics using the Particle-In-Cell (PIC) simulation code VPIC (Bowers et al., 2008). The properties of the spectral and PIC methods complement each other, so that such an approach allows for the most efficient exploration of the processes of interest.

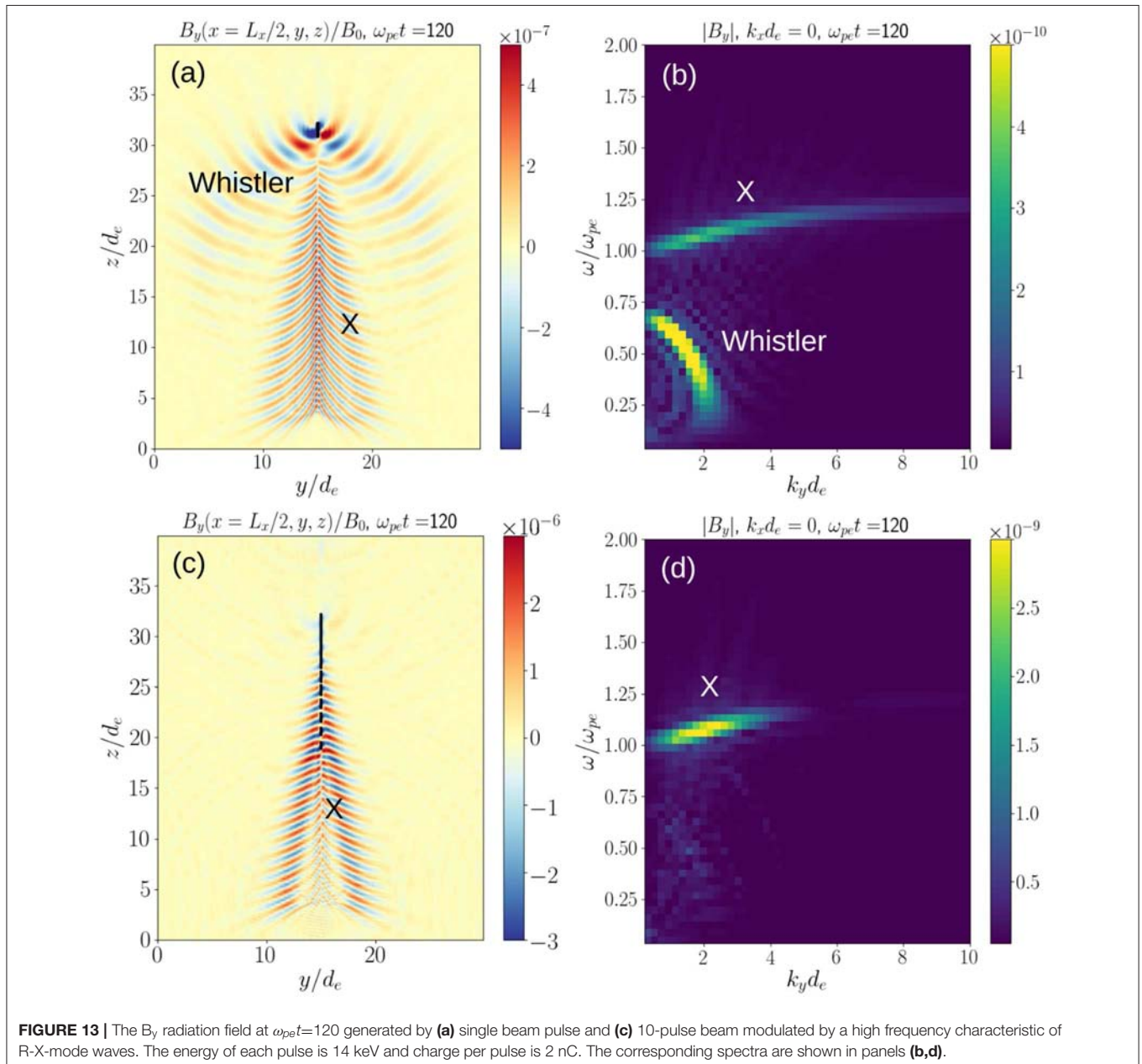
SPS uses a spectral decomposition of the plasma distribution function in terms of Hermite polynomials in velocity space, a Fourier decomposition in physical space (appropriate for problems with periodic boundary conditions) and an implicit time discretization. The velocity-space spectral decomposition is such that one can describe the plasma as a macroscopic fluid with the low-order moments of the expansion, while the kinetic physics is retained by adding more moments (Vencels et al., 2015). The beam simulations are performed using $N_H = 4$ Hermite polynomials in each velocity direction (corresponding to a fluid approximation) and the number of spatial modes along x , y , z directions equal correspondingly to 150, 150, and 248. The beam, implemented as an external current in the simulation, moves along z axis aligned with the external magnetic field B_0 .

Figure 13a shows the B_y component of the magnetic field (normalized to B_0) at the end of a simulation ($\omega_{pe}t = 120$) with $E = 14$ keV for a single beam pulse. The pulse is introduced into SPS as an external current produced by a cylindrical pulse with a total charge of 2 nC uniformly distributed inside the cylinder (shown as a black rectangle at $z/d_e \sim 32$). The initial



length of the pulse corresponds to a 100 ns beam pulse, while its radius is equal to $0.1 d_e$. In this SPS simulation, we follow the longitudinal dynamics of the pulse due to its own space charge but do not model the transverse dynamics of the pulse (see section First-principle simulations of Beam Dynamics). Also, we do not account for any potential neutralization of the beam pulse charge due to the background plasma. Two modes of radiation are clearly visible: the whistler mode (with longer wavelengths) and the R-X-mode (with shorter wavelengths). As the pulse propagates, its length increases and at $z/d_e \sim 25$ the length of the pulse becomes comparable to the wavelength of the R-X-mode plasma wave. Above this point the radiation is dominated by the whistler mode waves as radiation in the R-X-mode decreases

significantly due to coherence effects as expected from radiation theory (c.f. **Figure 12**). The wave spectrum of the radiation field from the single pulse simulation is shown in panel (b) where the presence of the two modes can be further differentiated. Panel (c) in **Figure 13** shows radiation from a 10-pulse finite-length beam with the same parameters as in **Figure 13a** but modulated by a frequency equal to $(\omega_{pe} + \omega_{uh})/2$ corresponding to a frequency of R-X-mode plasma waves. As expected from HB theory (**Figure 10**), the radiation is dominated by the R-X-mode. This is also confirmed by a spectrum of the radiation shown in panel (d). At $z/d_e \sim 27$ the pulses (shown by black rectangles) merge and the radiation field is dramatically reduced by coherence effects. **Figure 13c** suggests that, for the parameters



considered, radiation in the R-X mode would be maximized over $\sim 20 d_e$ (i.e., ~ 600 m). Future work will revisit this type of simulations including the transverse beam dynamics, beam energy spread and beam-pulse neutralization to make more accurate predictions of the extension of the R-X-mode wave field and optimize the separation between accelerator and receiver payloads accordingly.

A detailed comparison of wave generation and SPS simulations for a Gaussian pulse can be found in Delzanno and Roytershteyn (2019).

The type of simulations presented in **Figure 13** already remove some of the limitations of the HB theory. By accounting for finite-size beam pulses, thermal and non-linear effects, we can compute a finite radiated power in the whistler and R-X modes and compute the corresponding wave amplitudes for Beam PIE. In addition, a model for beam dynamics, which includes also the transverse dynamics, is being implemented in SPS, thus allowing calculations of the effect of the beam dynamics on the wave field and, in particular, of its extent before coherence effects reduce the waves amplitude.

First-Principle Simulations of Beam Dynamics

The spectral plasma solver (SPS) simulations described in the previous section allow us to assess properties of the (relatively weak) radiated field. However, their computational cost would increase dramatically if the beam dynamics and feedback between the beam pulses and the magnetized plasma were fully resolved. Particle in Cell (PIC) methodology offers a convenient alternative to study such processes. The trade-off is that the radiated field is not accurately described, mostly due to statistical noise associated with the finite number of computational particles.

Here we discuss some of the results from a preliminary VPIC simulation intended to study dynamics of beam pulses. The VPIC code solves a system of relativistic Maxwell-Vlasov equations for each plasma species in a 3D domain of spatial extent $L_x = L_y \approx 0.065 d_e$ and $L_z \approx 5.5 d_e$ with uniform magnetic field in the z direction. Boundary conditions corresponding to a perfect electric conductor are used on x and y boundaries and a perfect magnetic conductor on z boundaries. The number of cells is $200 \times 200 \times 16384$. The simulation is initialized with a uniform, two-component plasma with the following parameters: $\beta_e = \beta_i \approx 7 \times 10^{-7}$, $\omega_{pe}/\omega_{ce} = 1.3$, $m_i/m_e = 1,836$. As the simulation progresses, beam pulses with energy $E_b = 14$ KeV are injected at $z \approx 0$ with a δ -function distribution in energy. The beam electrons are treated as a separate plasma species. The initial beam radius is $r_b \approx 6.5 \times 10^{-4} d_e$ and the initial beam density is approximately 38 times higher than the background. The length of a single pulse is approximately $t_{p1}\omega_{pe} \approx 1$ and the time interval between pulses is $t_{p2}\omega_{pe} \approx 2\pi$. The beam is injected with zero pitch angle. The beam particles are absorbed at the other end of the domain at $z \approx L_z$. In the simulation discussed here, no additional positive charge is injected in the system to compensate for the injected beam charge.

As the beam pulse is injected, the electrostatic repulsion drives its rapid expansion transversely and, to a lesser

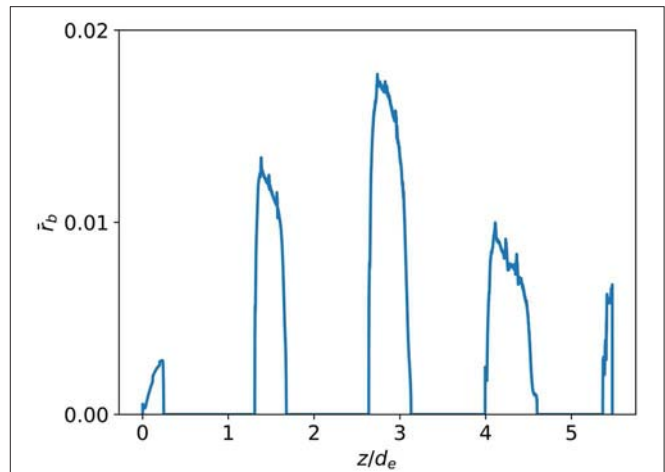
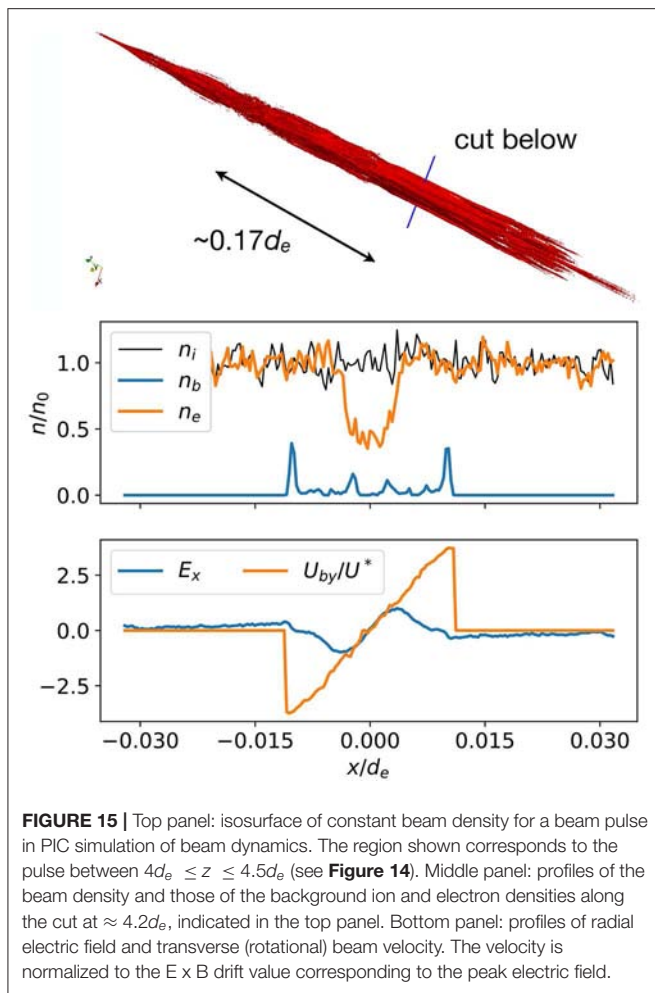


FIGURE 14 | Density-weighted radius (normalized to d_e) of beam pulses in a PIC simulation of beam dynamics as a function of the distance along the background magnetic field. The velocity is normalized to the $E \times B$ drift value corresponding to the peak electric field, while E_x is normalized to its peak value.

degree, longitudinally, along the beam propagation direction (also aligned with the background magnetic field). This is illustrated in **Figure 14**, which shows density-weighted mean-square radius of the pulses $r_b^2(z) = \int n_b(x, y, z) (x^2 + y^2) dx dy / \int n_b(x, y, z) dx dy$ as a function of z , the coordinate along the magnetic field. The transverse expansion is counterbalanced by the Lorentz force, such that mean beam radius oscillates in z . We note that in this simulation the pulses are injected with zero pitch angle spread and, hence, the dynamic value of the beam pulse radius is quite smaller than that used in the studies of **Figure 13**, $r_b = 0.1 d_e$. The transverse structure of a single pulse is shown in **Figure 15**. Here, the top panel shows an isosurface of constant beam density (beam density equal to 0.1 of the reference background density for the case shown). The middle panel shows profiles of background ion and electron densities, as well as the profile of the beam density along the cut indicated in the top panel. Profiles of the electric field E_x and the beam rotation velocity U_{by} along the same cut are shown in the bottom panel. Several important observations could be immediately made: i) individual beam pulses are shaped by a combined action of electrostatic repulsion, Lorentz forces, and instabilities; ii) the resulting *microscopic* structure is quite complex, but overall the beam pulse maintains coherence on spatial scales relevant to wave emission; iii) the beam pulse is partially charge-neutralized by the electrons of the background plasma, which reduces the severity of the electrostatic repulsion.

DISCUSSION

This paper describes the Beam-Plasma Interactions Experiment. Beam PIE is an active experiment that uses a novel linear accelerator based electron beam and advanced wave and plasma



diagnostics instruments to test beam wave-generation physics in greater detail than previously possible. The experiment uses a mother-daughter payload with the higher-altitude payload carrying the linear accelerator and the lower-altitude payload carrying a suite of wave and particle detectors. The first objective of the experiment is to conduct the first tests of modern linear accelerator technology in space which, if successful, could enable a new generation of active experiments. The second objective is to test the generation of whistler waves by the electron beam over a range of parameters (energy, modulation frequency, and duty cycle) not previously investigated in space. The third objective is to test the generation of R-X-mode radiation by the electron beam—an experiment that has not previously been done. The partitioning of energy and the detailed characteristics of the whistler vs. R-X-mode waves should be a sensitive test of our understanding of beam-plasma-wave interactions.

One example of a future active experiment is the Magnetosphere-Ionosphere Connections Explorer (CONNEX). The objective of this mission concept is to understand the magnetospheric processes that produce different types of auroral forms. It would use a multi-cell RF linear accelerator

to accelerate electrons to energies up to 1 MeV. The beam is strong enough to produce a visible spot in the auroral ionosphere to test the magnetic connectivity between the auroral ionosphere and the dipole-to-tail transition region in the equatorial magnetosphere. We note, however, that the design of the CONNEX beam would also allow pulsed-beam operations and therefore the opportunity to test the generation of waves by pulsed electron beams under magnetospheric (rather than ionospheric) conditions. Both CONNEX and the USAF DSX experiment (which uses a physical antenna to generate waves) could provide technology demonstrations for future active modification of the space environment including remediation of artificial radiation belts from High Altitude Nuclear Explosions (Reeves, 2018; Carlsten et al., 2019).

DATA AVAILABILITY STATEMENT

The datasets generated for this study are available on request to the corresponding author.

AUTHOR CONTRIBUTIONS

GR is PI of the Beam PIE mission and PF is deputy PI and payload manager. GD, KY, and VR are responsible for the simulation and modeling work shown. BC, JL, MH, and DN make up the electron accelerator team. RP is Goddard lead and responsible for the DC and AC electric field experiments as well as the fluxgate magnetometer. WF is responsible for the search coil. DR is responsible for the Langmuir probe. MS is responsible for the electron plasma spectrometer. ERS, ES, and ED are responsible for ground observations.

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Conflict of Interest: ERS was employed by company SRI International.

The remaining authors declare that the research was conducted in the absence of any commercial or financial relationships that could be construed as a potential conflict of interest.

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