

Electrode plasma transport into the gap of high-current Z-pinch accelerators

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(Received 14 November 2025; accepted 10 June 2026; published 1 September 2026)

In the pulsed-power accelerators that drive Z-pinch implosions, multi-MA currents rapidly heat the electrodes above the threshold for plasma formation. When these surface plasmas are transported into the transmission line gap, they shunt some portion of the driver current away from the load, referred to as “current loss.” In highly magnetized systems, current loss is well approximated by a Hall current proportional to the minimum plasma density in the gap and the electric field [N. Bennett, D. R. Welch, G. Laity, D. V. Rose, and M. E. Cuneo, Magnetized particle transport in multi-MA accelerators, *Phys. Rev. Accel. Beams*, 060401 (2021)]. This paper explores the mechanism that drives the initial plasma transport into the gap, which is difficult to isolate in large-scale power-flow simulations but is instrumental in seeding the current-loss mechanism. The stability and dynamics of an electron sheath emitted from a resistive cathode plasma in the presence of large magnetic fields are investigated using semianalytic theory and simulation. A dispersion analysis determines that the instability growth rate scales inversely to ion mass and magnetic field strength. This scaling is confirmed in kinetic models of an electron sheath and resistive plasma using the relativistic electromagnetic particle-in-cell code CHICAGO[®]. The simulations are initiated with field magnitudes representative of the final feed of a Z-pinch driver and demonstrate that a resistive plasma instability grows on ns timescales, generating perturbed electric fields of sufficient strength to rapidly accelerate plasma ions into the gap. Although the resulting plasma density in the gap is 2 orders of magnitude below the surface density, it is sufficient to generate the published current losses.

DOI: [10.1103/prpw-f8lt](https://doi.org/10.1103/prpw-f8lt)

I. INTRODUCTION

The multi-MA accelerators that drive Z-pinch implosions necessarily generate extreme current densities in the vicinity of the pinch. The electrodes carrying these high-current densities rapidly heat above the threshold for surface contaminant desorption and electrode plasma formation. These surface plasmas, when transported into the transmission line gaps, may shunt a portion of current before reaching the Z-pinch load, thus degrading load performance. This process, referred to as “current loss,” has been studied extensively for Sandia National Laboratories’ (>25 MA) Z accelerator [1–4].

Initial current-loss investigations focused on the larger radius, lower current-density regions of the accelerator,

such as the upstream magnetically insulated transmission lines (MITLs) and posthole convolute regions, in which magnetic nulls reside in complex geometries. There, the current density reaches 10^4 – 10^5 A/cm, and electrons, which are explosively emitted from the cathodes as in lower energy-density MITLs [5], are largely magnetically insulated and drift downstream. Some current loss occurs as nonmagnetically insulated particles cross the gaps near the magnetic nulls [6]. In the final transmission line feed to the Z-pinch load, or “inner MITL,” current densities in excess of an MA/cm produce magnetic fields exceeding MG levels, which rapidly Ohmically heat the electrodes. As the surface temperatures exceed 700 K, contaminants are liberated and surface plasmas quickly form in the presence of electron flow. Locally emitted electrons and electrode plasmas are highly magnetically insulated (the electron Larmor radius, r_L , is measured in μm for MG magnetic fields). Thus, the particle environment here is dominated by drifting electron sheaths and local surface plasmas. With this high magnetic insulation, current loss must proceed from conduction terms such as a Hall-related current loss [7] and instability-driven transport [8].

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A remaining open question is what mechanism drives the initial plasma transport into the transmission line gap. This physics is included, but difficult to isolate, in large-scale power-flow simulations. To address this question, this paper investigates the dynamics of electron sheath flow and surface plasma evolution and their role in transporting charged particles against the constraining magnetic field. A semianalytic theory of a resistive plasma instability (RPI) is compared to highly resolved particle-in-cell (PIC) simulations of the cathode electrode plasma and sheath using the CHICAGO[®] code [9]. The RPI described here follows the work of Swegle and Ott [10], who identified the sensitivity of MITL electron sheaths to destabilizing TM waves. Chernin [11], who extended this instability to lower frequencies by including electrode resistivity (resistive wall effect), describes the RPI as being driven by the drag force on the plasma electrons. Similarly, Chang *et al.* [12], who identified yet lower frequency modes and faster growth times by extending the resistive wall to include the more resistive electrode plasmas, identify the dissipation of the negative energy TM mode as the driver of the RPI. This work is not intended as a model of electrode emission turn-on, power flow, or current loss but analyzes the stability of evolving cathode plasmas in the unique field environment of the inner MITL of a Z-pinch driver. It provides insight into the physical mechanisms for rapid plasma transport that lead to current loss in the final inner MITL of the Z accelerator. Future work will include experimental validation of the electrode plasma density profiles.

Simulations demonstrate that the RPI leads to a rapid plasma density fill of the inner MITL. The dynamics that drive this instability are included, but obscured, in large-scale simulations of power flow and current loss. The instability is a result of the resistivity of the magnetized electrode plasma [13], as determined via dispersion analysis for linear growth in Sec. II. It is seen to grow inversely with the plasma ion mass and the applied magnetic field strength. This is confirmed in kinetic transmission line simulations presented in Sec. III. The simulations follow the RPI from initiation through the nonlinear regime, where plasma crosses the anode-cathode (AK) gap. Through RPI alone, plasma densities in the transmission line gap may exceed 10^{15} cm^{-3} , sufficient to conduct MAs of current across the inner MITL gap. Conclusions are discussed in Sec. IV.

II. STABILITY OF THE ELECTRON SHEATH EMITTED FROM A RESISTIVE PLASMA

As most commonly considered in MITL flow analysis, the dynamic evolution of the sheath electrons produces stable transverse profiles [14]. However, the sheared flow of an electron sheath emitted from a perfectly conducting boundary can lead to long wavelength, or diocotron, and short wavelength, or magnetron, instabilities [14].

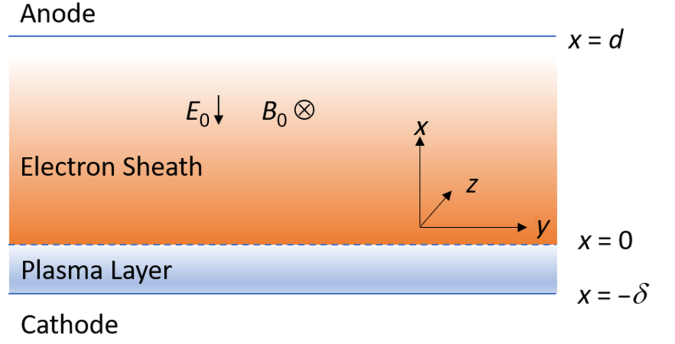


FIG. 1. Schematic of the planar sheath-plasma geometry assumed for a semianalytic analysis of the resistive plasma instability.

Furthermore, the coupling of the electron sheath to the plasma layer at the plasma-sheath boundary (configuration illustrated in Fig. 1) can lead to enhanced instability. We consider here the destabilizing effect of a thin surface cathode plasma layer on electron flow in an MITL. We generalize the previous linear analyses of the electron flow or resistive cathode instabilities [12–14] by including the ion contribution to the plasma conductivity tensor. A shooting method [15] is used to determine the dispersion relation. For fields and density profiles representative of a Z-pinch-driver inner MITL, robust instability occurs at short wavelengths and frequencies greater than the ion collision frequency. The calculated subnanosecond instability growth time is short compared with the typical 100-ns pulse power rise time.

A. Methodology

The equation of motion that defines the frequencies for the stability analysis assumes the geometry shown schematically in Fig. 1. The plasma extends from the metal (perfectly conducting) cathode at $x = -\delta$ to the electron sheath interface at $x = 0$. The equilibrium plasma density profile and temperature are prescribed and do not evolve in a stability analysis. The contribution of both electrons and ions in the immersed plasma is described by a tensor conductivity. The voltage and current are prescribed with an equilibrium consistent with a TEM-driven system. Following the practice of Refs. [12–14], and others, the equilibrium fields are perturbed by a TM wave that accommodates instability growth. The TM mode has a magnetic field purely transverse to the wave propagation direction with both transverse and longitudinal electric field components.

We first construct the conductivity tensor $\hat{\sigma}$ based on the generalized Ohm's law [16], where the diagonal elements are $\sigma_{\parallel} = \frac{nq^2}{m\nu}$ and $\sigma_{\perp} = \sigma_{\parallel}(1 + \frac{\omega_c^2}{\nu^2})$ for conduction parallel and perpendicular to the magnetic field, respectively, and $\sigma_H = \sigma_{\perp}(\frac{\omega_c}{\nu})$ is the magnitude of the off-diagonal elements, or the Hall conductivity. Here, n is the number

density of particles of mass m and charge q , with collision frequency and cyclotron frequency ω_c . For a collection of nonrelativistic charged particles (expected in the quasi-neutral thermal electrode plasma) with velocity $\bar{u}$, the equation of motion due to the temporal dependence of a perturbed TM mode of frequency ω ($e^{-i\omega t}$) is given by

$$-i\omega m\bar{u} = q\left(\bar{E} + \frac{\bar{u}}{c} \times \bar{B}_0\right) - \nu m\bar{u}, \quad (1)$$

where in this analysis $\bar{B}_0$ is the uniform, equilibrium magnetic field associated with the TEM mode. The tensor pressure and thermal force terms are neglected in Eq. (1) as impacting only isotropic thermal expansion. For motion in x and y in a z -directed magnetic field B_o , we obtain

$$(\nu - i\omega)mu_x = qE_x + q\frac{u_y}{c}B_o \quad (2)$$

and

$$(\nu - i\omega)mu_y = qE_y - q\frac{u_x}{c}B_o. \quad (3)$$

We construct the conductivity tensor for a collision frequency modified by the perturbed mode and ignoring conduction parallel to B_o . Taking the current density $\bar{J} = qn\bar{u}$ and inverting $\bar{J} = \bar{\sigma}\bar{E}$ gives

$$\bar{\sigma} = \frac{nq^2}{m\nu D} \begin{bmatrix} \nu'\nu & \frac{\nu q B_o}{mc} \\ \frac{-\nu q B_o}{mc} & \nu'\nu \end{bmatrix}, \quad (4)$$

where $\nu' = \nu - i\omega$ and $D = \nu'^2 + \omega_c^2$.

The electron conductivity matrix becomes

$$\bar{\sigma}_e = \frac{ne^2}{m_e\nu_e D_e} \begin{bmatrix} \nu'_e\nu_e & -\omega_{ce}\nu_e \\ \omega_{ce}\nu_e & \nu'_e\nu_e \end{bmatrix}, \quad (5)$$

where subscript e denotes electron, ν_e is the electron collision frequency with ions, $\omega_{ce} = \frac{eB_o}{m_e c}$ is the electron cyclotron frequency, $\nu'_e = \nu_e - i\omega$, and $D_e = \nu_e'^2 + \omega_{ce}^2$. Note that the tensor elements in Eq. (5) are similar to the generalized Ohm's law with a complex collision frequency.

The conductivity matrix for singly ionized ions is obtained from the electron tensor by switching from subscript e to i and reversing the signs of the off-diagonal Hall terms, i.e.,

$$\bar{\sigma}_i = \frac{ne^2}{m_i\nu_i D_i} \begin{bmatrix} \nu'_i\nu_i & \omega_{ci}\nu_i \\ -\omega_{ci}\nu_i & \nu'_i\nu_i \end{bmatrix}. \quad (6)$$

(For multiply-ionized ions, the charge in the coefficient is increased by the ionization state "Z.") We make use of

the total conductivity matrix $\bar{\sigma} = \bar{\sigma}_e + \bar{\sigma}_i$ in the calculations that follow.

For a description of the electron sheath in equilibrium, we use the laminar flow model described in Sec. III A of Ref. [14], which allows for softer transverse density profiles. The density profile was defined using dimensionless parameters as

$$N(s) = \frac{N_o \cos h(bs)}{1 + \exp\left[\frac{s-s_0}{a}\right]}, \quad (7)$$

where $s = x/d$ and d is the AK gap width. The specification of b , s_o , and a determines the overall profile shape, and the density N_o is chosen to give the desired gap voltage. The magnetic field B_o at the plasma-sheath interface fixes the impedance.

To investigate instabilities in the configuration in Fig. 1, we again follow the discussion in Ref. [14], assuming flute perturbations of wave number k on TM waves in the plasma and sheath of the form $e^{i(ky-\omega t)}$. We now allow that the longitudinal electric field at $x = 0$ will not, in general, be zero. This follows from the assumption of a resistive plasma from which the electrons are pulled and will determine the coupling of the sheath to the plasma. $E_y(0)$ is found by numerical integration of equations for perturbed fields E_x , E_y , and B_z in the plasma obtained from Ampere's and Faraday's laws:

$$\frac{\partial E_y}{\partial x} = ikE_x + i\frac{\omega}{c}B_z, \quad (8)$$

$$\frac{\partial B_z}{\partial x} = \frac{-4\pi}{c}(\sigma_{21}E_x + \sigma_{22}E_y) + i\frac{\omega}{c}E_y, \quad (9)$$

and

$$\left(\frac{4\pi}{c}\sigma_{11} - i\frac{\omega}{c}\right)E_x = ikB_z - \frac{4\pi}{c}\sigma_{12}E_y. \quad (10)$$

Note, σ_{ij} refers to the value of the total conductivity matrix in row i and column j . Equations (8)–(10) are integrated from $-\delta$ to 0 with initial conditions $E_y(-\delta) = 0$, $B_z(-\delta)$ arbitrary. At $x = 0$, $E_y(0)/B_z(0)$ gives the wave impedance $Z(0)$. This is matched to the wave impedance in the electron sheath at $x = 0$. For the integration across the sheath [17], we can take $B_z(0)$ to be arbitrary, $E_y(0) = Z(0)B_z(0)$, and use Eqs. (B-5), (B-6), (B-9), and (B-10) of Ref. [17] to solve for $\frac{\partial E_y}{\partial x}$ at $x = 0$.

These equations give the initial conditions for integration of the second-order differential equation in a non-linear, relativistic, equation shooting method described in Sec. III B of Ref. [14]. Shooting to satisfy the boundary condition $E_y(d) = 0$ for a given k determines the complex eigenvalue.

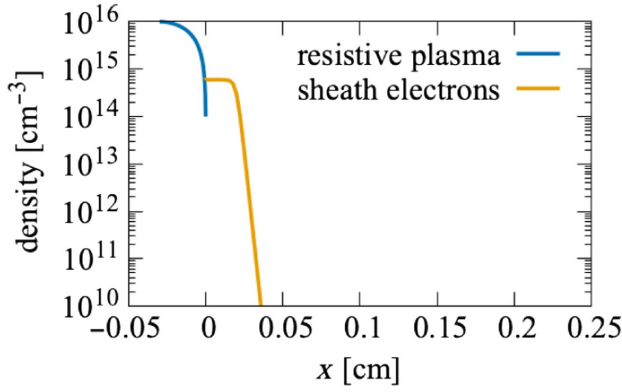


FIG. 2. Transverse density profile of the resistive plasma layer, including both electrons and ions, and the electron sheath density used for the analytic theory example are plotted. The sheath profile was taken from a simulation with perfectly conducting electrodes.

A summary of the solution steps is outlined here: (i) For a given k , make an initial guess for ω . (ii) Integrate the fields given in Eqs. (8)–(10) from $x = -\delta$ to 0 across the resistive plasma, with boundary condition $E_y(-\delta) = 0$, to obtain $E_y(0) \wedge B_z(0)$. (iii) Using the transverse fields at $x = 0$ from (ii), integrate across the sheath to obtain $E_y(d)$ (using Eqs. (B-5), (B-6), (B-9), and (B-10) of Ref. [17]). (iv) Making use of a numerical nonlinear solver, adjust the complex frequency ω and repeat (i)–(iii) until the condition $E_y(d) = 0$ is met to the specified accuracy. Of course, $\omega_i > 0$ implies instability.

B. Numerical results

We solve the system of equations with boundary conditions described above for the eigenmodes in an MITL with a 4.5-MV anode potential, a 570-kG anode magnetic field, and an AK gap $d = 0.25$ cm. The density profile of the electron sheath is taken from an idealized simulation with the above gap, potential, and current, and perfectly conducting electrodes for $x \leq 0$ instead of a plasma. Space-charge-limited emission of electrons is permitted from the cathode to produce the sheath. Note that for this idealized simulation using a perfectly conducting boundary, the growth rates are extremely small due to the rounded electron sheath profiles. Only by adding the resistive or dissipative collisional plasma as assumed in the analytic theory above, where $E_y(0) \neq 0$, do we significantly destabilize the sheath. The electron sheath profile from the simulation was fit using Eq. (7), with $a = 0.0065$, $s_o = 0.0907$, $b = 2.0$, and $N_o = 4.5810^{14}$ cm⁻³ to give the correct gap voltage. The density profiles for the resistive plasma layer approximate those observed in MITL PIC simulations with evolving plasma density (see, for example, Fig. 4 in Ref. [13]). We assume a characteristic plasma layer of thickness $\delta = 0.03$ cm with a density of 10^{16} cm⁻³ at the cathode, dropping quadratically

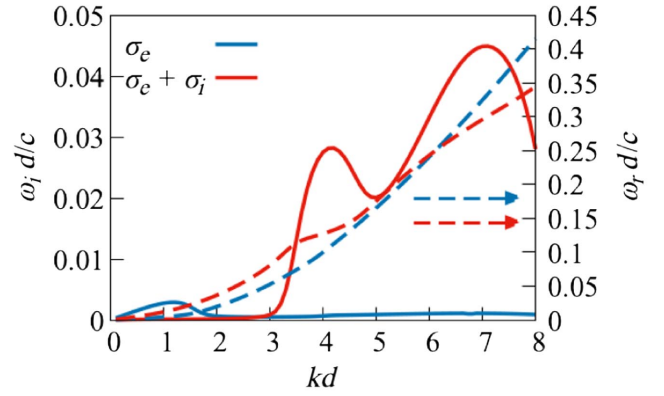


FIG. 3. Normalized growth rates of the perturbed TM wave for mobile electrons only in the plasma layer (case 1 labeled σ_e in blue) and for mobile electrons and ions (case 3 labeled $\sigma_e + \sigma_i$ in red), and the imaginary (solid lines) and real frequencies (dashed lines) are plotted versus wave number.

to 10^{14} cm⁻³ at the plasma-sheath interface, as shown in Fig. 2 along with the sheath electron distribution from Eq. (7). The plasma temperature is a uniform 3 eV. The results below are not sensitive to plasma temperature until the collision rate approaches the cyclotron frequency, which occurs at a much colder 0.1-eV temperature for 10^{16} cm⁻³ density. Spitzer collision rates [18] are calculated from the density and temperature and are used in the calculation of $\bar{\sigma}$.

The instability growth rates for the perturbed TM mode are calculated for four cases: (1) including only the electron contributions to $\bar{\sigma}$, (2) including only the electron contributions and only the off-diagonal terms in Eq. (5), (3) including both electron and ion contributions to $\bar{\sigma}$, and (4) including both electron with off-diagonal terms in Eq. (5) and ion contributions with only the diagonal terms in Eq. (6). These four cases were chosen to elucidate the new effect of the ions in the conductivity matrix. Also, the cases show how the electron component is largely due to the cross magnetic field or Hall conduction because the electrons are magnetized, and the ion component is largely due to perpendicular conduction to B_0 because the ions are not magnetized.

The results are shown in Fig. 3 for the full-conductivity cases (cases 1 and 3). Note that the electron-only peak growth occurs near wavelength $\lambda = 1.27$ cm ($kd = 1.24$) with e-folding time $t_e = 2.82$ ns ($\omega_i d/c = 0.00296$). For case 3, the peak growth rate is an order of magnitude higher (0.16 ns e-fold time) and occurs at a higher k with 0.22-cm wavelength and two peaks. In both cases, the real frequencies ($\omega_r d/c$) rise with k . Case 1 has a 1.38-ns real oscillation period at peak growth, while case 3 has a 0.0279-ns period at peak growth.

We now compare the dispersion calculation using the complete conductivity matrices (cases 1 and 3) to those using the aforementioned approximate electron and ion conductivity descriptions (cases 2 and 4) in Fig. 4. The

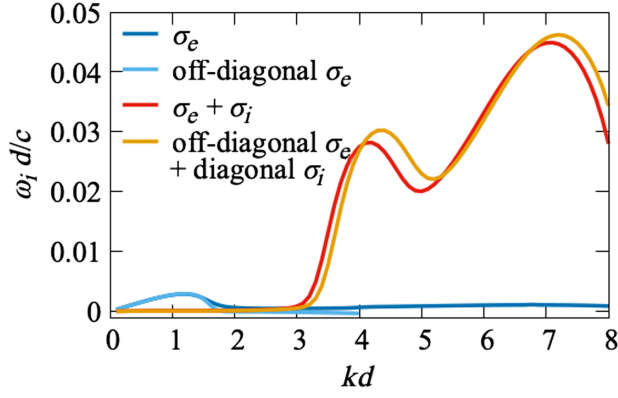


FIG. 4. Normalized growth rate of the perturbed TM wave versus wave number for mobile electrons only in plasma layer (case 1 in blue), approximate electron-only growth rate with only off-diagonal components (case 2 in cyan), growth rate for mobile electrons and ions (case 3 in red), and approximate growth rate with off-diagonal electron and diagonal ion contributions to tensor conductivity (case 4 in yellow).

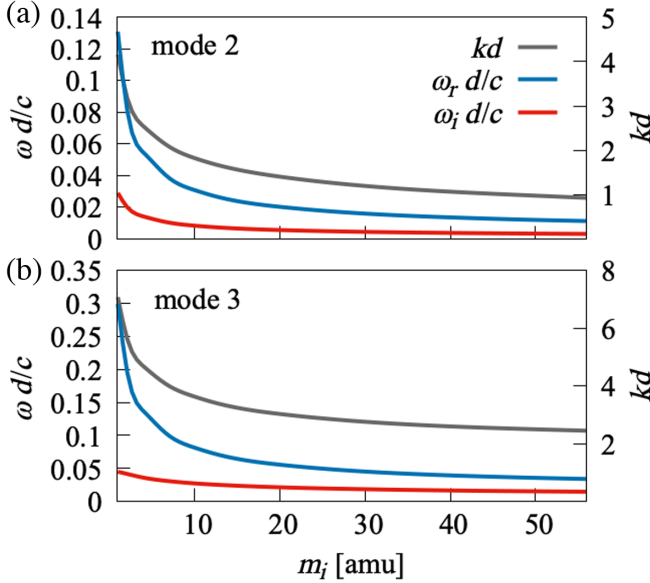


FIG. 5. Normalized wave number and instability growth rates, kd , ω_r , and ω_i , calculated using the full-conductivity matrix (full $\sigma_e + \sigma_i$) at 570 kG, are plotted as functions of ion mass for (a) mode 2 and (b) mode 3.

growth rates produced from cases 2 and 4 are excellent approximations for the peak growth rate from cases 1 and 3 for the chosen plasma density profiles. We can neglect the off-diagonal ion contributions to $\bar{\sigma}_i$ since ω_{ci} is relatively small given the large ion mass. The assumption of neglecting electron diagonal terms is valid for lower plasma densities but can break down at higher densities due to increased electron collisionality. If we reduce the minimum plasma layer density at the plasma-sheath interface to $< 10^{14} \text{ cm}^{-3}$, little change in theoretical growth rate

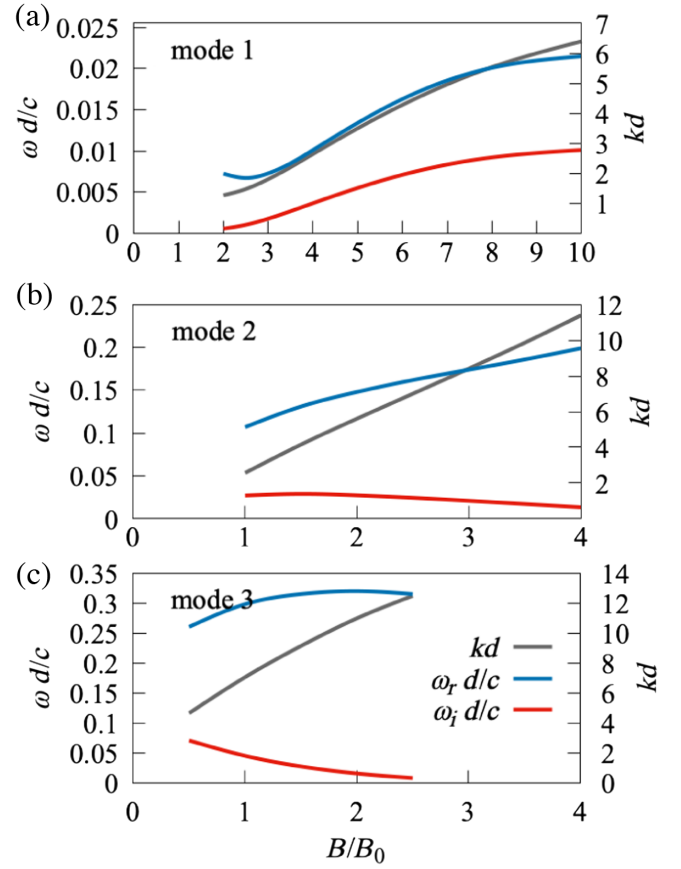


FIG. 6. Normalized wavelength and instability growth rates, kd , ω_r , and ω_i , calculated using the full-conductivity matrix with a proton mass, are plotted as functions of a normalized applied equilibrium B field. The values are shown for (a) mode 1 (full σ_e), (b) mode 2 (full $\sigma_e + \sigma_i$), and (c) mode 3 (off-diagonal σ_e + diagonal σ_i).

behavior is calculated; however, increasing the minimum density to $> 2 \times 10^{15} \text{ cm}^{-3}$, the growth rate decreases by an order of magnitude.

1. Ion mass and magnetic field dependence

Although the density profile shown in Fig. 2 is valid for a nominal ion mass (1 amu) and baseline magnetic field (570 kG), we now make use of this profile while we vary these parameters to get an approximate sensitivity to m_i and B . We can make reasonable arguments for the ion mass dependence by assuming the density distribution evolves more slowly with higher mass. We expect the peak growth rate to fall with mass and approach the electron-only growth at very large mass. We find that the two peaks observed in Fig. 4 persist at large m_i . We refer to these as ion modes 2 and 3 because of the inclusion of the ion tensor conductivity. We refer to the case with the full electron tensor conductivity and no ion conductivity as mode 1. The values for k and ω for modes 2 and 3 are plotted in Fig. 5. Mode 1 is neglected in Fig. 5 because it

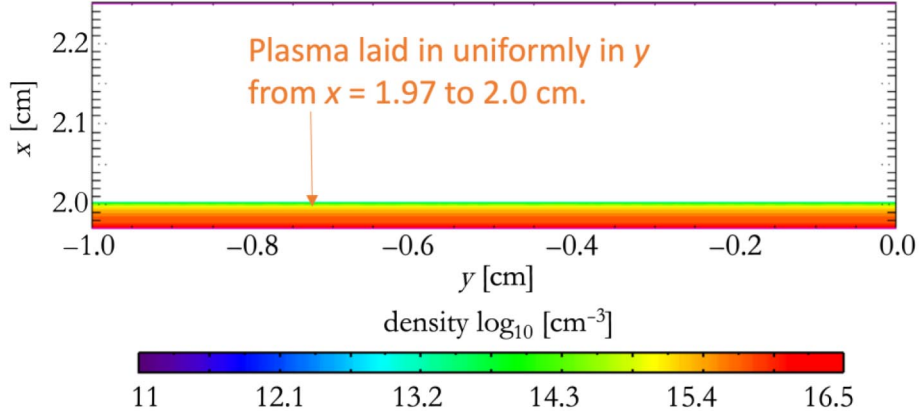


FIG. 7. The 2D Cartesian simulation setup for studying the RPI. Plasma is initialized on the cathode with a uniform 570 kG B_z field and 4.5-MV potential drop across the anode and cathode.

was overwhelmed by the other modes and not visible with the shooting method. Using the full electron and ion tensors (full $\sigma_e + \sigma_i$), the growth rate for mode 3 exceeds that of mode 2 for all masses, given the baseline 570-kG B field. The growth rate falls more slowly than linearly with m_i . For singly ionized carbon, the rate is roughly half that of a proton.

Although the magnetic field will modify the sheath profile, we examine the dispersion behavior by varying the field about the baseline B_0 value. We can track the variation of the three individual modes by slowly varying the ratio B/B_0 as shown in Fig. 6. We find that mode 1 asymptotes to the electron-only value at large m_i , and we will refer to it as the electron mode. Modes 2 and 3 are only seen with finite ion mass and can be referred to as ion modes. Unlike the electron mode, these modes have a decreasing growth rate with B . Mode 2 decreases more gradually than mode 3 and dominates the growth at $B/B_0 > 2.5$. This is due to the decreased diagonal contribution with increasing ion cyclotron frequency. In all modes, the real frequency and wave number generally increase with B . The two ion modes do not show significant instability for $B/B_0 > 4$.

III. CHICAGO SIMULATIONS OF RESISTIVE PLASMAS WITH SHEATHS

The semianalytic theory in Sec. II is tested using a kinetic transmission line simulation of the Fig. 1 geometry. The analytic theory can only examine small-amplitude wave growth of the plasma and electron sheath system. In addition to this linear behavior, the kinetic theory also examines the local evolution of the sheath and plasma densities and temperatures. This evolution can lead to vortices, electric field reversal, and plasma transport into the gap. The simulations do not include electrode emission turn-on, impedance transition, power flow from upstream, or current loss. A planar transmission line is modeled in 2D Cartesian coordinates using the fully relativistic,

electromagnetic PIC code CHICAGO. A fully kinetic plasma is placed on top of a perfectly conducting cathode to a depth of 0.03 cm, as shown in Fig. 7. To isolate the plasma/sheath stability, a baseline simulation is initialized with this plasma, periodic boundaries in y , an applied 4.5-MV voltage, and uniform 570-kG magnetic field. Note that the plasma surface is now at $x = 2$ cm and the anode surface is at $x = 2.25$ cm, not at $x = 0$ and 0.25 cm as in the theory section. The kinetic simulation includes binary Coulomb scattering between all particles as described in Ref. [13]. The fields are initialized with an electrostatic solution (using a 1D Poisson solve $\partial E_x / \partial x = \rho$) and a constant applied magnet field. The particles and

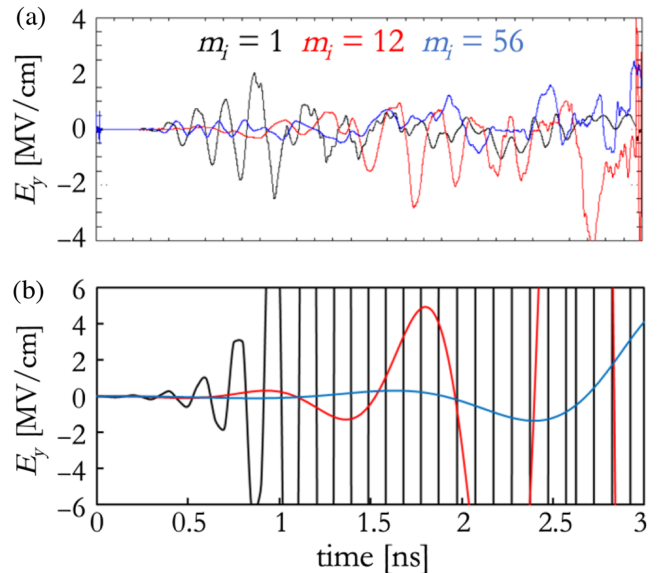


FIG. 8. (a) Simulated longitudinal electric field (E_y) just within the electrode plasma at $y = -0.1$ cm and $x = 2$ cm is plotted for ion masses of H^+ , C^+ , and Fe^+ ($m_i = 1, 12$, and 56) with 570-kG magnetic field. (b) E_y from the analytically calculated highest growth rate for the same masses. The linear theoretical growth continues unbounded.

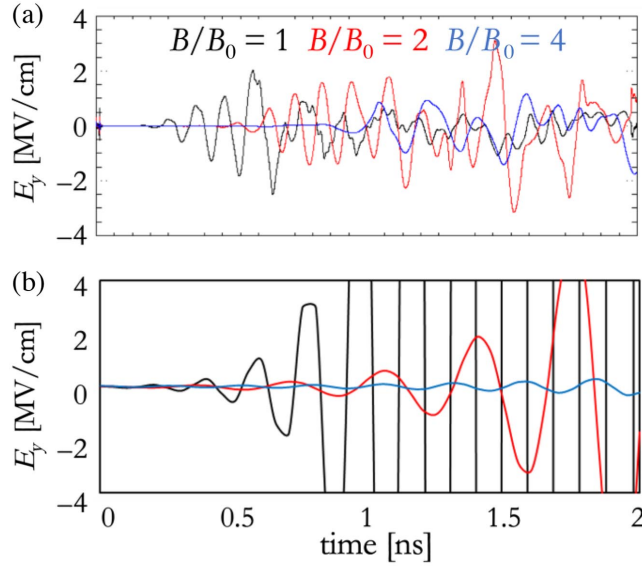


FIG. 9. (a) Simulated longitudinal electric field (E_y) just within the electrode plasma at $y = -0.1$ cm and $x = 1.99$ cm is plotted for field strengths of $B/B_0 = 1, 2$, and 4 with the ion mass of H^+ . (b) E_y from the analytically calculated highest growth rate for the same field strengths. The linear theoretical growth continues unbounded.

electromagnetic fields are then advanced with an updated Magnetic Implicit technique [19]. The simulations demonstrate good numerical convergence with $\Delta x = 0.00125$ cm and $\Delta y = 0.0025$ cm spatial resolution

(2–4 cells per collisionless plasma skin depth) and 0.04-ps time step (peak plasma frequency-time step product, $\omega_{pe} \Delta t = 0.25$). We used 900 electrons and 225 ions as initial macroparticles per cell. Energy conservation over the 2–5 ns simulation duration was better than 0.5%.

Charged particles can be transported from the surface of the plasma, so neutrality is restored by enabling emission from the underlying solid cathode. To enable this, the electrode becomes a space-charge-limited electron emitter after 240 kV/cm electric field threshold is exceeded. Longitudinal electric field probes record the field values near the plasma outer surface ($x = 2$ cm). This calculation only crudely approximates conditions relevant to Sandia's Z accelerator inner MITL (for example, in Refs. [7,13]) since the pulse rise is neglected and, with it, its impact on plasma evolution.

To best compare with the semianalytic theory, the y extent of the simulation corresponds to the wavelength determined in theory to be the most unstable. The longitudinal electric field (E_y) is plotted in Fig. 8 at $y = -0.1$ cm and $x = 2$ cm for a series of simulations with plasma ion masses corresponding to H^+ , C^+ , and Fe^+ . Figure 8 also shows E_y from the highest analytically calculated growth rate assuming a 20 kV/cm perturbation. The real frequency at peak growth from theory, shown in Fig. 5, falls with mass from 4 to 1 GHz as the growth e-fold time ($1/\omega_i$) increases from 0.14 to 0.56 ns. Comparing the simulated fields [Fig. 8(a)] with those of the RPI theory [Fig. 8(b)], we see the simulated growth rate resembles

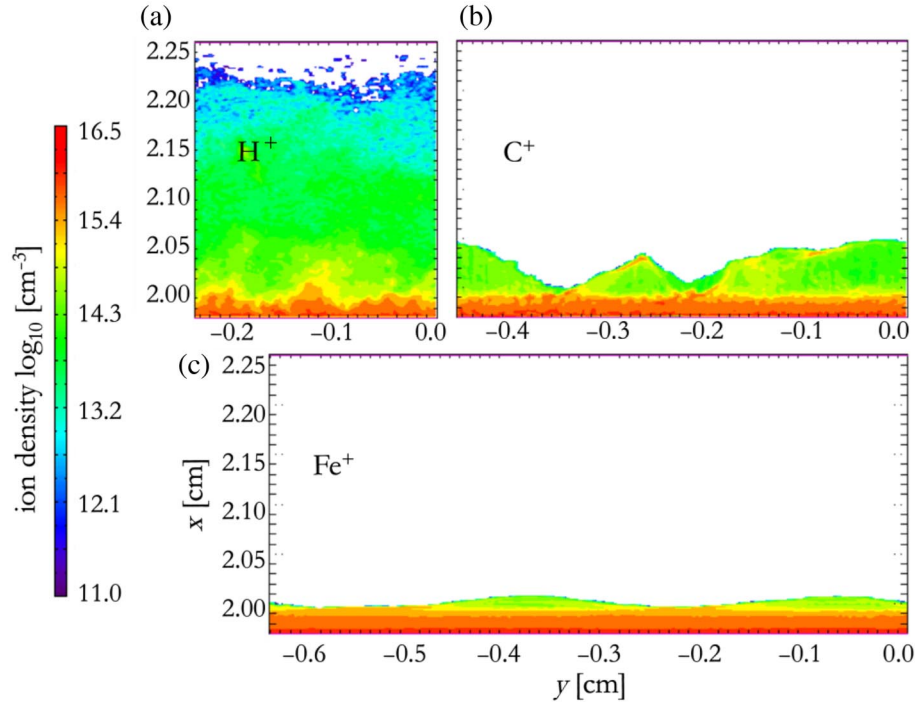


FIG. 10. Plasma ion density at $t = 2$ ns is plotted for (a) H^+ , (b) C^+ , and (c) Fe^+ simulations at 570 kG. Note that the simulated wavelength increased from 0.22 cm for H^+ to 0.65 cm for Fe^+ .

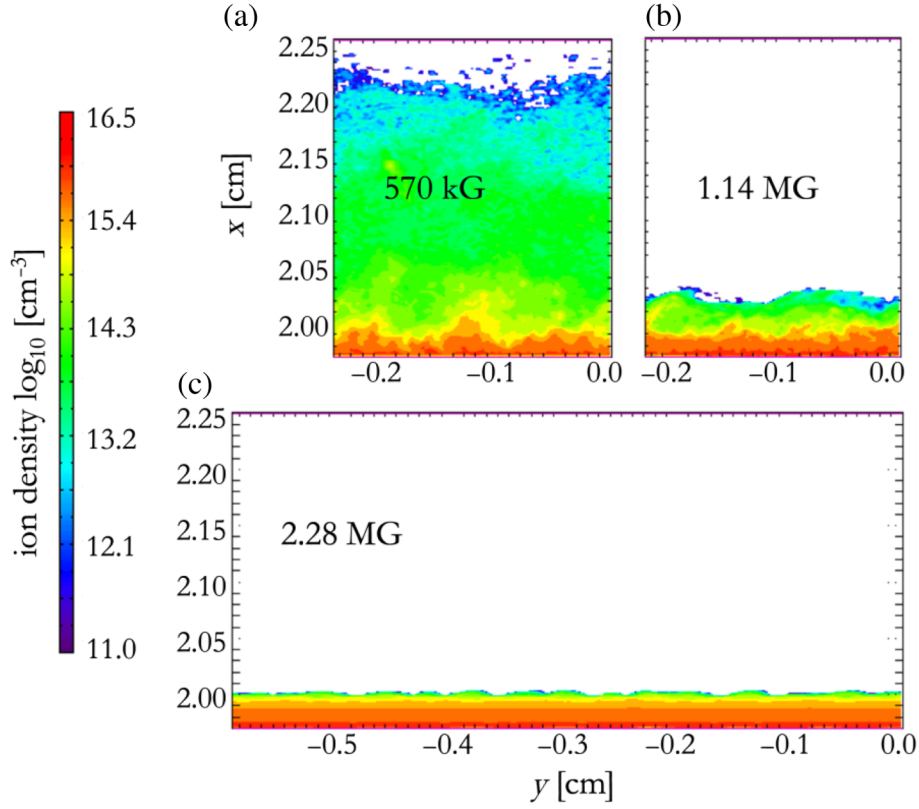


FIG. 11. Plasma ion density at $t = 2$ ns is shown for magnetic fields corresponding to $B/B_0 =$ (a) 1, (b) 2, and (c) 4 for $B_0 = 570$ kG and fixed ion mass (H^+). The simulated wavelength increases from 0.22 to 0.59 cm to accommodate the fastest growing mode.

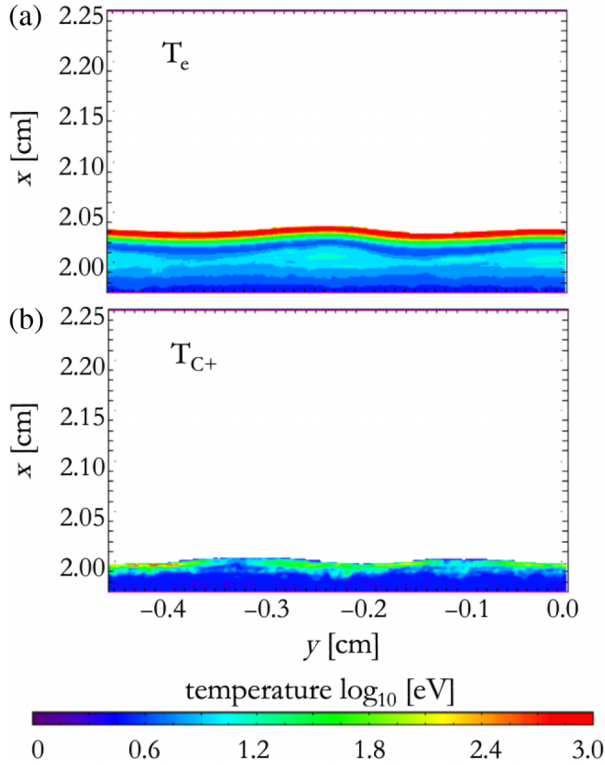


FIG. 12. At $t = 1.2$ ns, the electron (a) and C^+ (b) temperatures are plotted for the e- C^+ plasma at 570 kG.

those from theory until E_y reaches 2000 kV/cm, where the instability saturates from disassembly of the plasma. The linear theoretical growth continues unbounded. For the larger ion mass, we do see the predicted lower frequency modes as well as coupling from higher frequency modes. As with the theory, the growth rate is proportional to approximately $1/\sqrt{m_i}$.

Figure 9 plots the E_y instability growth, as in Fig. 8, but for H^+ ions with varying magnetic fields of $B/B_0 = 1, 2$, and 4. Again, the simulation length in y is taken to be that of the wavelength of the mode with the largest instability growth from the analytic theory. We expect the analytically calculated growth rates to be roughly that of the simulations since both theory and simulation assume the same electron sheath transverse profile for all values of B/B_0 .

The simulated plasma ion densities for the three ion masses at $t = 2$ ns are plotted in Fig. 10. As expected, the progression of ions across the gap is slower for higher mass. The density within the gap, well off the surface, is order 10^{14} cm^{-3} for the lighter two ions. Figure 11 shows the H^+ plasma densities for $B/B_0 = 1, 2$, and 4, also at 2 ns. The plasma transport is slower for higher magnetic fields. While the H^+ ions exhibit rapid transport for $B/B_0 = 1$ and 2, reaching 10^{14} cm^{-3} density in the gap, the $B/B_0 = 4$ case has only just entered the nonlinear phase by 2 ns. This transport behavior occurs well into the

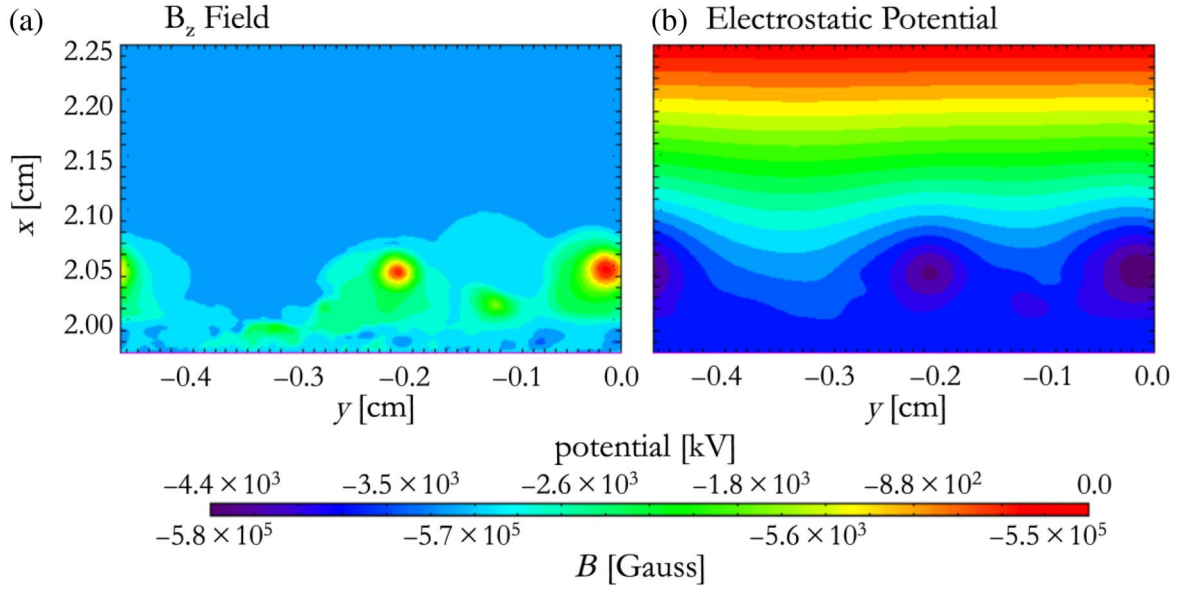


FIG. 13. Magnetic fields are plotted in (a) for the C⁺ mass and 570-kG simulation. In (b), electrostatic potentials are plotted. The plots are at 2 ns.

nonlinear phase of the instability when ions are reacting to the highly perturbed electron sheath flow.

The nonlinear process by which the plasma ions are transported into the AK gap can be deduced from the

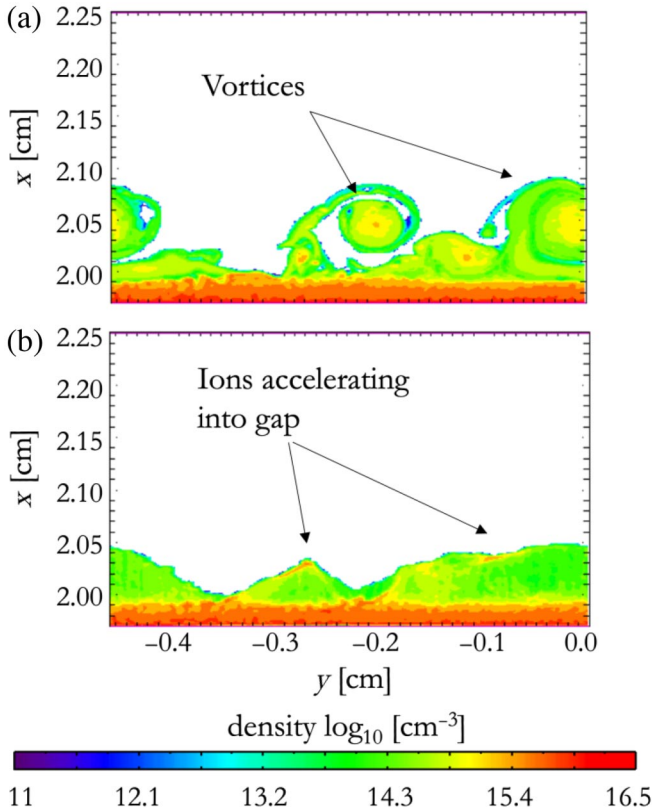


FIG. 14. Electron (a) and C⁺ (b) densities are plotted at 2 ns. The transport is driven by MV-scale potential islands.

simulations. The large longitudinal electric field and sheath oscillations near the resistive plasma surface destabilize the plasma electron sheath. As shown in Fig. 12, the surface waves heat both the electrons and ions to nearly keV temperatures by 1.2 ns, reducing the plasma collisionality further. High space-charge electron vortices develop (see Fig. 13) reminiscent of those caused by abrupt transitions in transmission line impedance [20]. The vortices cause the electric field to reverse as electrostatic islands form with potentials on the order of MV. Thus, the charge neutralization of the ions is severely degraded. The vortices affect both the electric and magnetic fields as shown in Fig. 13. We see at 2 ns two vortices, shown clearly in Fig. 14, produce sufficient electric field to accelerate the ions toward the anode. By this time, the ions have moved 0.5 mm into the gap.

IV. SUMMARY AND CONCLUSIONS

With semianalytic theory and CHICAGO simulations, we have studied the growth of the RPI and its nonlinear impact on plasma transport into an AK gap. The theory, with sheath and plasma parameters informed by simulation, is solved with a shooting method to obtain the dispersion relation. We find that including ions in the conductivity matrix has a large impact on the RPI growth. The ion contributions dominate the diagonal components, with the electrons dominating the off-diagonal components. The general trends show an increase in RPI e-fold growth time with ion mass and magnetic field. These trends are backed up with detailed simulations of particle transport in a 0.57 to 2.28-MG transmission line.

The CHICAGO simulations also follow the impact of the instability on the plasma and how plasma is transported into the gap. The strong longitudinal oscillations quickly destabilize the plasma electrons, resulting in vortices due to wave breaking. These vortices sustain MV-scale potentials to which the ions respond and move into the gap. This work provides a good physics explanation for previously published simulation work that showed rapid plasma motion into the gap, resulting in current loss for the Z accelerator. In Ref. [7], 10^{15}-cm^{-3} plasma densities cross the AK gap in an MG-scale magnetic field, similar to the conditions in Figs. 11(b) and 11(c), generate up to 3 MA of current loss along a roughly 3-cm-long transmission line segment.

The mechanism driving the instability is described in Refs. [11] and [12], which assert that coupling of a cathode electron sheath with collisional plasma layers leads to a plasma drag force or dissipation of the negative energy TM mode. This particle drag reinforces the drag force, thus driving instability. In our analysis of a highly magnetized plasma, the plasma electrons exhibit only weak dissipation at very low frequency, which is inhibited by the strong confining magnetic fields. Including the new ion contributions to the conductivity matrix extends and enhances the instability at higher ω_r because these ions are not magnetized and exhibit more dissipation. The RPI enables the growth of an electric field component parallel to the sheath drift. This perturbs the plasma surface, generating flow with vortices that enable plasma transport into the gap against a constraining magnetic field.

The implications of this work are that significantly reducing the inventory of lighter surface contaminants, chiefly water, would delay the transport of plasma into the gap. The motion of heavier substrate ions, such as iron, would take roughly the square root of the mass longer (assuming the substrate ions are singly ionized). A further implication is that the transport of plasma into the gap scales roughly with $1/B$ or is Bohm-like. Thus, delaying the production of electrode plasma until the magnetic fields are higher would help; however, the Ohmic heating and plasma formation closely follow the magnetic field strength. In general, this behavior bodes well for current transport in higher current facilities. We hope to quantify and incorporate this research into faster, more accurate models of accelerator performance in the future.

ACKNOWLEDGMENTS

Lawrence Livermore National Laboratory is a multi-mission laboratory managed and operated by Lawrence Livermore National Security, LLC under prime contract DE-AC52-07NA27344 for the U.S. Department of Energy's National Nuclear Security Administration. Sandia National Laboratories is a multi-mission laboratory managed and operated by National Technology & Engineering Solutions of Sandia, LLC, a wholly owned

subsidiary of Honeywell International Inc., for the U.S. Department of Energy's National Nuclear Security Administration under contract DE-NA0003525. We also appreciate thoughtful discussions with David Rose of Pacific Fusion.

DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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