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Do dispersive waves play a role in collisionless magnetic reconnection?

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Using fully kinetic simulations, we demonstrate that the properly normalized reconnection rate is fast ~0.1 for guide fields up to 80× larger than the reconnecting field and is insensitive to both the system size and the ion to electron mass ratio. These results challenge conventional explanations of reconnection based on fast dispersive waves, which are completely absent for sufficiently strong guide fields. In this regime, the thickness of the diffusion layer is set predominantly by the electron inertial length with an inner sublayer that is controlled by finite gyro-radius effects. As the Alfvén velocity becomes relativistic for very strong guide fields, the displacement current becomes important and strong deviations from charge neutrality occur, resulting in the build-up of intense electric fields which absorb a portion of the magnetic energy release. Over longer time scales, secondary magnetic islands are generated near the active x-line while an electron inertial scale Kelvin-Helmholtz instability is driven within the outflow. These secondary instabilities give rise to time variations in the reconnection rate but do not alter the average value. © 2014 AIP Publishing LLC.

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I. INTRODUCTION

Magnetic reconnection often leads to explosive release of energy and is thought to be operative in space, laboratory, and astrophysical plasmas.1,2 Observational constraints require the rate to be fast—i.e., independent of the system size and operating on Alfvénic time scales. Simulations in the collisionless regime have shown the reconnection rate to be \( u_{in}/V_A \sim 0.1 \) over a range of parameters in the limit of zero guide field3–6 and in the presence of a guide field.7–10 Here, \( u_{in} \) is the ion inflow velocity and \( V_A \) is the ion Alfvén speed. However, one of the outstanding issues is the identification of the mechanism responsible for these fast rates. In steady state, the reconnection rate is proportional to the aspect ratio of the diffusion region, which develops structures on both ion and electron kinetic scales. While the thickness of these structures can be understood from simple scaling arguments, the physics responsible for regulating the length of the diffusion region remains controversial.11

One prominent idea has been that kinetic reconnection is crucially dependent on the presence of fast dispersive waves,12–14 including whistler and/or kinetic-Alfvén waves (KAW) depending on the specific regime.15 Within this line of reasoning, the presence of waves with quadratic dispersion (\( \omega \propto k^2 \)) below the ion-scale causes the electron outflow velocity to scale inversely with the width of the layer (\( u_{e,\text{out}} \sim \omega/k \propto k \sim 1/\delta \)). This allows the mass flux (\( \delta u_{e,\text{out}} \)) to remain constant even as the layer thickness \( \delta \) becomes very small, leading to a rate that is insensitive to the mechanism which breaks the frozen-in condition.

While these intuitive arguments are appealing, there are two lines of evidence that challenge this interpretation. First, the same fast reconnection rate ~0.1 also occurs in electron-positron plasmas16–18 which do not support any fast dispersive waves.19 Although it has been argued that temperature anisotropy instabilities may provide an alternate explanation,20 fast rates are still observed when these instabilities are suppressed.21 Second, hybrid simulations in which the Hall term is removed have also reported fast reconnection22 in the absence of dispersive waves. However, the interpretations of these results remain controversial,23 and thus there is interest in finding alternative ways to rigorously isolate the influence of dispersive waves.

With this goal in mind, we re-examine the influence of dispersive waves by using the guide field to change the properties of the waves permitted in the system. As shown from two-fluid theory,15 no quadratic waves exist in the strongly magnetized limit defined by \( \beta = 8\pi n(T_e + T_i)/B_0^2 < m_e/m_i \) and \( B_0^2 \ll B_T^2 m_e/m_i \) where \( B_T \equiv B_0(1 + b_x^2)^{1/2} \) is the total magnetic field, \( B_0 \) is the reconnecting component, \( b_x \equiv B_x/B_0 \) where \( B_x \) is the guide field and \( m_e/m_i \) is the electron to ion mass ratio. In this regime, fast dispersive waves are eliminated and the remaining waves have negative dispersion at small scale. Previous two-fluid simulations have shown a clear transition to slow reconnection in this regime.15 For less extreme conditions, both two-fluid simulations24–27 and laboratory experiments28 suggest a transition to faster reconnection rates when the thickness of the diffusion region falls below the ion sound radius \( \rho_s \). However, a variety of scaling dependencies on \( \rho_s \) have been reported in fluid simulations,29–33 while a recent fluid-kinetic model34 in the regime \( \beta \sim m_e/m_i \) reported fast rates comparable with the weak guide field limit.

One of the issues with the two-fluid studies is that they often exhibit formation of nearly singular sublayers,31,35–38 with thickness limited by the imposed dissipation. To avoid this issue, here we use fully kinetic particle-in-cell
simulations which provide the most complete treatment of collisionless plasmas. In contrast to two-fluid simulations, our kinetic simulations suggest these sublayers are limited by finite electron gyro-radius effects. In addition, we demonstrate that the properly normalized reconnection rate remains nearly invariant $\sim 0.1$ for strong guide fields $\beta \ll m_i/m_e$ where fast dispersive waves are suppressed. These fast rates are insensitive to the system size, the mass ratio, and the initial layer thickness, suggesting that dispersive waves are not pertinent. Furthermore, these simulations reveal interesting new results concerning the dynamics of reconnection in the limit of ultra strong guide fields $b_g = 80$, which are much larger than previously considered $b_g \sim 4$ in kinetic studies.\textsuperscript{7–10} These results are potentially relevant to studies of strong magnetohydrodynamic (MHD) turbulence where the current sheets formed tend to possess strong guide fields due to the strongly anisotropic cascade to small scales.\textsuperscript{36} In addition, these regimes are relevant to astrophysical applications, where the increasing importance of the displacement current limits the Alfvén velocity to the speed of light, while order unity deviations from charge neutrality emerge within highly extended electron layers. As a result, a significant portion of the energy released from the reconnected flux goes toward building up intense electric fields in the outflow. Finally, both secondary tearing and a chain of Kelvin-Helmholtz (KH) vortices are observed in these elongated layers, resulting in time variability in the reconnection rate, but not altering the average value.

This paper is organized as follows. First, the basic simulation setup employed for this study is described in Sec. II. Next in Sec. III, a series of simulations are considered in which the guide field is varied $b_g = 1 \rightarrow 80$ along with the mass ratio $m_i/m_e = 1 \rightarrow 25$ and the system size. To permit the broadest range of guide fields, the initial layer half-thickness was set equal to the electron inertial length. To examine the influence of this choice, we consider a simulation with an ion-scale initial layer in Sec. IV with a guide field of $b_g = 8$, and demonstrate that our conclusions from Sec. III remain valid in these broader current sheets. Finally, in Sec. V, we summarize our findings and speculate on a possible explanation for these results.

II. SIMULATION SETUP

The fully kinetic two-dimensional simulations in this study were performed using the VPIC code.\textsuperscript{40} For the series of simulations discussed in Sec. III, the initial condition is a force-free current layer with $\mathbf{B} = B_0 \tan(z/\lambda) \mathbf{x} + B_0 \hat{b}_g^2 \hat{y}$ + sech$^2(z/\lambda)$, corresponding to a field of magnitude $B_T \equiv (B_0 + \hat{b}_g^2)^{1/2}$ which rotates by an angle $\phi = 2\tan^{-1}(1/b_g)$ across a layer with half-thickness $\lambda$. The initial particle distributions are Maxwellian with spatially uniform density $n_0$ and temperature $(T_i = T_e)$. Ions are initially stationary while electrons have a net drift $\mathbf{U}_e$ to produce a current density $\mathbf{J} = -e n_0 \mathbf{U}_e$ consistent with $\nabla \times \mathbf{B} = 4\pi \mathbf{J}/c$. Spatial scales are normalized by the inertial length $l_i = e\omega_{pi}/c$ for each species $(j = i,e)$, where $\omega_{pj} = (4\pi n_0 e^2/m_j)^{1/2}$ and time is normalized by the cyclotron frequency using the reconnecting field $\Omega_j = eB_0/(m_j c)$. The boundary conditions are periodic in the $x$-direction, while on the $z$-boundaries particles are reflected and the field boundary conditions are conducting. While force-free current layers are generally considered to be more realistic in low-$\beta$ regimes, the conclusions drawn in this paper also apply to the standard Harris sheet\textsuperscript{43} equilibrium with a strong uniform guide field, as demonstrated in Sec. IV.

III. INFLUENCE OF DISPERSIVE WAVES

To examine the influence of dispersive waves, a series of runs are considered with guide field ranging from $b_g = 1 \rightarrow 80$ and mass ratio $m_i/m_e = 1 \rightarrow 25$. The electron thermal velocity is held fixed $v_{te}/c = \sqrt{T_e/m_e} = 0.1$ along with $\omega_{pe}/\Omega_{ce} = 5$. For this series of runs, the initial current sheet thickness is $\lambda = d_0$ which permits a reasonably fast onset of reconnection over this wide range of guide fields. However, we note that the results of this study only apply on a-shape $(\lambda = d_0)$ initial layers as demonstrated in Sec. IV. Unless otherwise noted, in this section the simulation domain size is $L_x \times L_z = 25.6d_0 \times 6.4d_0$, which is comparable with that used in Ref. 15. This corresponds to $n_e \times n_i = 4096 \times 6144$ cells with 150 particles per cell in the major run discussed (case 3). A localized perturbation is used to initiate reconnection in a controlled manner.

Parameter regimes that support quadratic waves ($\omega \propto k^2$) below ion scales are denoted as fast dispersive, while regimes with negative dispersion (i.e., shorter wavelength modes propagate slower) are referred to as slow dispersive. While not the main focus of this paper, the $m_i/m_e = 1$ limit is a nice reference since there are no fast dispersive waves.\textsuperscript{19} On the other hand, the $m_i/m_e = 25$ simulations permit a continuous transition between the fast and slow dispersive regimes, based on two governing parameters, $\mu_g = (1 + \beta/2)(1 + b_g^2)(m_i/m_e)$ and $\beta_0 = 8\pi n_0 (T_i + T_e)/B_0^2$ which is held fixed at $\beta_0 = 1$ in this study. For reference, the relevant regimes are illustrated graphically in Fig. 2 of Ref. 15. In this paper, our primary focus is on three cases: (a) Case 1 with $b_g = 1$ and $\mu_g \approx 0.1$ with fast dispersion on the whistler branch, (b) Case 2 ($b_g = 10$ and $\mu_g \approx 4$) only permits waves with slow dispersion, and (c) Case 3 ($b_g = 80$ and $\mu_g \approx 256$) is in the extreme limit of slow dispersion, where the Alfvén velocity for parallel propagating waves is relativistic.

To confirm the dispersive properties for these three cases, the phase velocity computed from linear Vlasov theory (solid lines) is shown for the whistler branch in Fig. 1(a) and for the KAW branch in Fig. 1(b). For each case, the propagation angle relative to the magnetic field $\theta = \tan^{-1}(b_g)$ is held fixed to examine the in-plane waves that are permitted in the 2D simulations. Only case 1 supports fast dispersion on the whistler branch while all three cases have slow dispersion for the KAW branch. For comparison in case 1, we also include the two-fluid dispersion relation with electron inertia\textsuperscript{15} as well as the Hall MHD limit ($m_e \rightarrow 0$) where the whistler dispersion supports fast dispersive waves all the way down to the dissipation scale, usually set by hyper-resistivity or the grid scale. In contrast, the inclusion of electron inertia causes the waves to slow down well before $kd_e < 0.5$ as the whistler transitions into electron
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FIG. 1. Dispersion relation from linear Vlasov theory for whistler (a) and KAW waves (b) for \( b_x = 1 \) (case 1), \( b_x = 10 \) (case 2), and \( b_x = 80 \) (case 3). The blue dashed line is from two-fluid theory for case 1, while the blue dotted-dashed line in (a) corresponds to the Hall MHD \( m_i = 0 \) limit. For each case, the angle with respect to the magnetic field is \( \theta = \tan^{-1}(b_x) \), \( V_A \) is the Alfvén speed, and \( V_{Ae} = V_A \cos \theta \). Panel (c) shows the normalized reconnection rate \( R \) as function of guide field \( b_x \) for a fixed system size \( L_x/d_x = 128 \), while panel (d) shows the reconnection rate as function of system size for fixed guide field \( b_x = 80 \).

cyclotron waves. Notice this happens even sooner in the Vlasov treatment, presumably due to kinetic damping.

While the wave dispersion changes dramatically between these cases, the normalized reconnection rate is nearly constant between \( b_x = 1 \rightarrow 80 \) as shown in Fig. 1(c). In each of these simulations, the reconnection rate is computed from

\[
R = \frac{1}{B_0 V_{Ae}} \left\langle \frac{\partial \psi}{\partial t} \right\rangle,
\]

where \( \psi = \max(A_y) - \min(A_y) \) along \( z = 0 \), \( A_y \) is the y-component of the vector potential, \( \langle \cdot \rangle \) represent a time average over \( \Delta t \Omega_L = 2 \) to reduce noise, and \( V_{Ae} \) is the Alfvén wave velocity based on reconnecting field \( B_0 \)

\[
V_{Ae} = V_A \cos \theta = \frac{v_A \cos \theta}{\sqrt{1 + (v_A/c)^2}}, \quad (1)
\]

where \( v_A = B_T/\sqrt{4\pi n(m_i + m_e)} \) is the non-relativistic Alfvén speed based on the total magnetic field. The values reported in Fig. 1(c) correspond to the late-time evolution after a quasi-steady reconnection process has been set up. Cases 1 and 2 are in the non-relativistic regime \( (v_A < c) \) so that \( V_{Ae} \approx v_A \cos \theta \), while case 3 is in the relativistic limit \( (v_A/c \approx 3.2) \) which implies \( V_{Ae}/c \approx \cos \theta \). The generalized expression in Eq. (1) results from including the displacement current in the linear theory for Alfvén waves. As shown later in this section, this same limiting outflow velocity can be derived from a Sweet-Parker type analysis, including deviations from charge neutrality. The reconnection rates for the \( m_i = m_e \) simulations are in good agreement with \( m_i/m_e = 25 \) in Fig. 1(c), suggesting a close link in the essential physics in these two limits. Finally, Fig. 1(d) shows that the reconnection rate for case 3 \( (b_x = 80) \) remains independent of system size, in stark contrast to the slow reconnection predicted for this regime by two-fluid simulations.

Next, we consider the dynamical evolution of the \( b_x = 80 \) simulation (case 3) in more detail. The initial current layer thickness \( \lambda = d_0 \) is much larger than the gyroradius for both the ions \( (\lambda/\rho_i \approx 23) \) and the electrons \( (\lambda/\rho_e \approx 114) \). To examine the kinetic physics in this regime, the time evolution of the reconnection rate for case 3 is shown in Fig. 2(a), while the structure of the reconnection layer at late time is shown in Fig. 3. The evolution can be roughly divided into two phases: (a) an onset period \( \Delta t_{de} < 20 \) in which the time derivative electron inertia, \( m_e \partial_t (n_e u_e) \), is the only non-ideal term available to break the frozen-in condition and (b) a quasi-steady phase \( \Delta t_{de} > 20 \) where terms arising from electron pressure tensor \( \nabla \cdot \mathbf{P}_e \) become dominant near the x-line. The duration of the onset phase is determined by balancing the electron inertial term against the reconnection electric field at the x-line (Fig. 2(b)), which implies the electron fluid velocity near the x-line increases linearly \( \Delta u_{ef} \approx e\Delta A_y/(m_e c) \) since \( n_e \) remains constant. This leads to the formation of an intense current sublayer with half-thickness \( (2 - 3)\rho_e \) near the x-line (Fig. 2(d)), consistent with previous kinetic studies. However, even at late time, most of the integrated current for the diffusion region is spread over a broader \( d_e \)-scale layer as illustrated in Fig. 2(d) and suggested in Refs. 38 and 42. To examine this issue in more detail, the spatial extent of the non-ideal \( E_e + (u_e \times \mathbf{B})_y \neq 0 \) region is shown in Fig. 4. The white box approximately encloses the region where \( E_e + (u_e \times \mathbf{B})_y > 0 \) with the exception of the sublayer which remains strongly non-ideal in the downstream region. As indicated by the black electron streamlines, the electron inflow is diverted in the outflow direction beginning near the \( d_e \) scale. Along the outflow boundary, the mass flux features a strong peak with half-thickness \( \sim 5\rho_e \) corresponding to the electron sublayer (see shaded red region and inset). However, this thin layer accounts for only ~20% of the total mass flux, while 80% passes through the broader \( d_e \)-scale layer.

The various terms in the generalized Ohm’s law are shown in Fig. 2(c) across the layer for the onset phase \( (t_1) \), and in Fig. 2(d) for the fully developed nonlinear phase \( (t_2) \). In each case, the results are time averaged over an interval \( \Delta t = 1/\Omega_L \) to reduce noise. As anticipated above, electron inertia is the dominant non-ideal term during the onset phase, with \( m_e \partial_t (n_e u_e) \) balancing the reconnection electric field at
the \(x\)-line. At later times, the formation of an embedded \(q_e\)-scale current sheet gives rise to a non-gyrotropic electron pressure tensor which produces \(r/C_1 P_e\) terms near the \(x\)-line. Although these new terms are crucial for setting up a quasi-steady process, notice that the reconnection rate in Fig. 2(a) does not change substantially between \(t_1\) and \(t_2\).

Finally, at even later times \(t \gg \Omega ci/30\), oscillations in the rate occur in conjunction with secondary island formation in the diffusion region current sheet (Fig. 5(a)), however, the average reconnection rate is not substantially altered.

As expected, the ion outflow in Fig. 3(b) approaches the Alfvén velocity \(V_{Ai}/c \approx \cos \theta \approx 0.012\) (Fig. 3(b)). The outflow velocity in the limit of very strong guide field can be derived from the following analysis. Including displacement current and charge separation, the momentum equation is \(\partial [n \rho_e \mathbf{u} + \mathbf{E} \times \mathbf{B}/4 \pi c] / \partial t + \nabla \cdot (n \rho_e \mathbf{u} u_{\|}) = 0\), where \(\mathbf{T} = (\mathbf{B} \mathbf{B}^T/2 + \mathbf{E} \mathbf{E}^T/2)/\epsilon_0\) is the Maxwell stress tensor and \(\mathbf{P} \approx \rho_e \mathbf{I} + (\mathbf{P}_{\perp} - \mathbf{P}_{\|}) \mathbf{y} \mathbf{y}\) is the pressure tensor for a strong guide field in the \(y\)-direction. Note that for cases studied here, the plasma flows are non-relativistic. Following a similar procedure to that of Ref. 48, we consider

![FIG. 2. (a) Time evolution of normalized-reconnection-rate (i.e., from measuring the change of global magnetic flux) for case 3. In (b), the correlation of \(\Delta \Theta_p\) and \(\Delta \Lambda_r\) shows onset phase where electron inertia is dominant and the quasi-steady phase where \(\nabla \cdot P_e\) terms become important. The various terms in the time-averaged momentum equation (normalized by \(n_0 V_{Ai} B_0/c\)) are shown in (c) at \(t_1 = 15/\Omega ci\) and in (d) at \(t_2 = 30/\Omega ci\), along with the current density \(J_y\).](http://example.com/fig2)

![FIG. 3. Results for case 3 with \(b_y = 80\) during the fully developed quasi-steady phase \((t_2 = 30/\Omega ci)\) showing (a) current density \(J_y\), (b) ion outflow velocity \(u_{iz}\) with cut at \(z = 0\) over-plotted with \(u_{ix}\), (c) electric field \(E_z\) with cut at \(z = 0\) over-plotted with \(-u_{iz} B_y/c\), and (d) charge separation \((n_e - n_i)/n_0\). White lines in (a), (c), (d) are flux surfaces and black lines in (b) are electron streamlines.](http://example.com/fig3)

![FIG. 4. (a) The non-ideal electric field of case 3 at \(t_2 = 30/\Omega ci\) with superposed electron streamlines shown in black curves. The white box corresponds to the approximate region where \(E_z + (u_e \times B)/c > 0\) and is the same region depicted in Fig. 3(b). (b) The electron mass flux is shown along the right boundary of the white box (\(x = 10 d_e\)). The shaded red region in the sublayer corresponds to only 20% of the mass flux through the layer. The inset shows the structure of this inner sublayer which is limited by finite gyroradius effects.](http://example.com/fig4)
incompressibility is assumed. Note that deduced from the frozen-in condition gyroradii are small compared with the gradient scales. Thus, velocity of both species is close to the E-vective (i.e., sublayer where the electron outflow speed exceeds $v_A$) with the exception of the narrow ion-scale sheet. Instead of the force-free setup in Sec.II, this section considers. In this section, we consider a larger system with an doubly periodic boundary conditions. Thus, the initial magnetic field is given by $E/V_A$ and using $P_{\perp,\text{out}} - P_{\perp,\text{in}} \sim B_{\perp,\text{in}}^{2}/8\pi$ while along outflow boundary this gives $P_{\perp,\text{out}} - P_{\perp,\text{in}} \sim B_{\perp,\text{in}}^{2}/8\pi + E_{\perp,\text{in}}^{2}/8\pi$. Hence, $P_{\perp,\text{out}} - P_{\perp,\text{in}} \sim E_{\perp,\text{in}}^{2}/8\pi$ and using Eq. (2), we obtain the outflow velocity

$$u_{\text{out}} \sim \frac{v_A \cos \theta}{\sqrt{1 + (v_A/c)^2}}$$

consistent with Eq. (1) and the previously derived expression for reconnection outflow speed in relativistic regimes.

The in-plane electron flow velocity remains very close to the ions throughout the domain, with only small differences within the $\rho_\perp$-scale current sheet (see comparison cut below Fig. 3(b)). This feature is dramatically different from all previous simulations of collisionless reconnection, where the electrons strongly decouple from the ions, resulting in outflow speeds that scale with the electron Alfvén velocity. In contrast, for these strong guide field regimes, the in-plane velocity of both species is close to the $E \times B$ drift, since the gyroradii are small compared with the gradient scales. Thus, the streamlines for the electron flow (black lines in Fig. 3(b)) follow the ion flow, with the exception of the narrow $\rho_\parallel$-scale sublayer where the electron outflow speed exceeds $V_A$, but remains well below the electron Alfvén velocity.

As shown in Fig. 3(c), the reconnection outflow induces intense electric fields associated with the nearly ideal convection (i.e., $E_z \sim -u_{\text{out}} B_z/c$), with a number of interesting implications. First, the $E_z$ field is associated with order unity deviations from charge neutrality as per Gauss’s law $E_z/\delta \sim 4\pi e(n_i - n_e)$. Since the layer thickness is $\delta \sim 0.5\lambda_i$, this implies $(n_i - n_e)/n_0 \sim 2h_0(\Omega_{ce}/\omega_p)(u_{\text{out}}/c) \sim 0.4$, which is consistent with Fig. 3(d). Since the electron gyroradius is much less than the Debye length ($\rho_e/\lambda_\parallel \approx 0.06$), these strong deviations from charge neutrality persist, despite the presence of Kelvin-Helmholtz instability (Fig. 5(b)).

Second, the $E_z$ field is important in the force-balance across the layer $\partial(P_{\perp,\text{out}} + B^2/8\pi - E_{\perp,\text{in}}^2/8\pi)/\partial z \sim 0$, and in setting the limiting outflow in Eq. (3). Although the reconnection rate is insensitive to the guide field, particle energization is dramatically reduced for regimes with $v_A > c$. This property follows immediately from ideal convection with outflow speed $u_{\text{out}}/c = \cos \theta$, which implies the energy density in the outflow electric field is comparable with the inflow reconnecting magnetic field $E_{\perp,\text{in}}^2/8\pi \approx B_{\perp,\text{in}}^2/8\pi$, thus absorbing a significant fraction of the total energy release.

Over longer time scales, the extended diffusion region current sheet illustrated in Fig. 3 is unstable to both secondary magnetic islands and a KH type instability as illustrated in Fig. 5. The formation of secondary islands gives rise to variations in the reconnection rate (see Fig. 2(a)) consistent with previous kinetic simulations. However, over the available duration in this simulation, the average reconnection rate appears unchanged. The presence of secondary Kelvin-Helmholtz instabilities in reconnection outflows has been predicted theoretically within a MHD description, but to our knowledge has never been observed in MHD simulations. However, the KH vortices in Fig. 5(b) are on the $d_\perp$-scale, where the MHD predictions do not apply directly. Previous kinetic simulations with much weaker guide fields have reported that electron-scale KH can give rise to the formation of secondary islands. In the present simulations, we found no clear evidence of this effect. In particular, notice that the flux surfaces in Fig. 5(b) are not wrapping with the vortices, indicating the frozen-in condition is violated within these $d_\perp$-scale structures and thus no magnetic islands are generated. In contrast, the secondary islands observed in the simulation occur near the existing x-line, where the in-plane velocity shear is weak but magnetic shear is strong, suggesting these islands are the result of a tearing-type instability.

IV. ION-SCALE INITIAL LAYER

In Sec. III, the initial layer half-thickness was equal to the electron inertial length $\lambda = d_i$ for all the simulations considered. In this section, we consider a larger system with an ion-scale sheet. Instead of the force-free setup in Sec. II, this simulation used the standard Harris sheet equilibrium with doubly periodic boundary conditions. Thus, the initial magnetic field is given by

$$B(z) = B_0 \left[ \tanh \left( \frac{z + L_z/4}{\lambda} \right) - \tanh \left( \frac{z - L_z/4}{\lambda} \right) \right] \hat{x} + B_y \hat{y},$$

where $B_0$ is the asymptotic field in the $x$-direction, $B_y = 8B_0$ is a uniform guide field in the $y$-direction, and $\lambda = d_i$ is the half-thickness of the layer. The initial density
is provided by drifting Maxwellian distributions with spatially uniform drift velocity $U_{ij} = 2cT_j/(q_iB_0d)$ and uniform initial temperature $T_i = T_e$ for both species. In addition, a uniform non-drifting background population is included with density $n_b = n_s$ and the same temperature. The system size is $L_x \times L_z = 128 \times 32 \times d$, and is resolved by $n_s \times n_b = 8192 \times 2048$ cells, using mass ratio $m_i/m_s = 25$, equal temperatures $T_i = T_e$ and $\nu_{pe}/\Omega_{ce} = 2.5$ (defined using $n_s$ and $B_0$), which implies $v_{th/c} = 0.2$. To initiate reconnection in a controlled manner, a magnetic perturbation of the form

$$
\delta B_x = -\frac{\delta B}{2} \left( \frac{L_x}{L_z} \right) \sin \left( \frac{2\pi x}{L_x} \right) \sin \left( \frac{4\pi z}{L_z} \right),
$$

$$
\delta B_z = \frac{\delta B}{2} \cos \left( \frac{2\pi x}{L_x} \right) \sin^2 \left( \frac{2\pi z}{L_z} \right)
$$

is imposed with $\delta B = 0.03B_0$. The relevant dispersion parameters ($\mu_i = 2.62$ and $\beta_i = 1$) are close to case 2 in Sec. III, and thus only slow dispersive waves are supported.

Since this case is doubly periodic, there are two separate current sheets with very similar behavior. For simplicity, we focus on one of these layers. The resulting reconnection rate is shown in Fig. 6 along with the current density at the select times indicated. The onset of reconnection is considerably slower in this thicker layer and proceeds in two steps. First, the overall current sheet is compressed by the long wavelength perturbation as illustrated in Fig. 6(b). During this phase, there is very little reconnection. Next, within this compressed region, a chain of tearing islands form and trigger fast reconnection ($R = 0.06$) as the islands grow to large amplitudes (see Fig. 6(c)). As reconnection proceeds, these initial magnetic islands are swept into the downstream (see Fig. 6(d)), while new secondary islands are formed within the $d_s$-scale diffusion region current sheet as shown in Fig. 6(e). Modest variations in the reconnection rate are observed as secondary islands form and limit the length of the diffusion region current sheet. However, the average reconnection rate does not change substantially over the duration considered. Despite the absence of fast dispersive waves, all of the above features are entirely consistent with previous kinetic simulations in weak guide field regimes. Furthermore, the late-time evolution of this ion-scale layer is in good agreement with all the results presented in Sec. III for electron-scale initial layers. In particular, the bulk of the current within the diffusion region is supported by a $d_s$-scale layer, while a smaller fraction supported by a $\rho_i$-scale sublayer and the resulting reconnection rate is nearly same.

Thus while the onset of reconnection in $d_s$-scale layers takes longer, the subsequent nonlinear evolution is very similar to the $d_s$-scale layers discussed in Sec. III.

V. SUMMARY

The fully kinetic results in this paper demonstrate that the normalized reconnection rate is insensitive to the guide field or to the presence of fast dispersive waves, and thus offer a clear test both for future theoretical models and reduced simulation approaches (either fluid or kinetic). In particular, recent gyrokinetic simulations have demonstrated fast reconnection rates in this regime which are in good agreement with the predictions of this manuscript. Nevertheless, without a rigorous quantitative theory of collisionless reconnection, there remains some ambiguity in the interpretation of these results. One might argue that fast dispersive waves are crucial in some regimes, but that other presently unidentified mechanisms are active for pair-plasma reconnection, or for the strong guide field regimes in this manuscript. However, to seriously argue for this scenario would require some remarkable coincidences.

First, the new mechanism would have to turn on at the same point in parameter space where the fast dispersive waves are suppressed. After all, if this new mechanism was active in the standard regime, then there would be two simultaneous explanations. In this case, which mechanism is really dominant? Second, if the reconnection physics in these regimes is fundamentally different, one would naturally expect some modulation in the reconnection rate and/or non-ideal dissipation terms across the transition in parameter space. As shown in this manuscript, both the reconnection rate and dominant terms in the Ohm’s law remain unchanged as fast dispersive waves are gradually suppressed.
While still speculative, one possible explanation for this persistence of fast reconnection in these very different parameter regimes is that the most unstable collisionless tearing modes occur for $k_0 \sim 0.5$ within kinetic scale current layers,56,57 regardless of the mass ratio or guide field. Assuming that the most unstable tearing modes constrain the length of the diffusion region, the condition $k_0 \sim 0.5$ implies an aspect ratio of $\delta/(2\pi/k) \sim 1/4\pi \sim 0.08$, which is consistent with the results of this paper. In particular, the large ion-scale simulation shown in Fig. 6 supports this general idea, both during the onset phase as fast reconnection first develops and also during the late-time evolution as secondary islands emerge. This idea may be connected with recent fluid theory and simulations,42 where it was suggested there is a critical value of the tearing stability index (i.e., $\Delta$ parameter) necessary for fast reconnection in collisionless regimes and the fastest growing modes also occur for $k_0 \sim 0.5$. However, a clear and convincing resolution to these questions will require further work. One important question, which is beyond the scope of this paper, is why the previous two-fluid simulations obtained a slow rate,15 in these strongly magnetized regimes. We are currently exploring a number of ideas and will report on these in a future publication.

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