Fundamental role of ion viscosity on fast magnetic reconnection in large-guide-field regimes

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Nonlinear analytical theory of magnetic reconnection with a large guide field is presented for the first time. We confirm that two distinct steady-state reconnection regimes are possible depending on the relative size of the diffusion region thickness δ versus the sound gyroradius ρ_s. The reconnection is slow (Sweet–Parker-like) for $\delta \gtrsim \rho_s$, and fast otherwise. However, unlike earlier work, we find that ion viscosity μ plays a fundamental role in the fast regime. In particular, for $\delta < \rho_s$ we obtain $\delta \propto \mathrm{Ha}^{-1}$, with $\mathrm{Ha} \propto 1/\sqrt{\eta\mu}$ as the Hartmann number, and the reconnection rate $E_z \propto \mathrm{Pr}^{-1/2}$, with $\mathrm{Pr} = \mu/\eta$ as the Prandtl number and η as the resistivity. If the perpendicular ion viscosity is employed for μ, the reconnection rate becomes independent of plasma β and collision frequencies, and therefore potentially fast. © *2010 American Institute of Physics*. [doi:10.1063/1.3449589]

Fast magnetic reconnection phenomena observed in space and laboratory plasmas indicate that nonlinear reconnection rates are independent of collisional dissipation coefficients and system size. Such dissipation independent rates have been reproduced computationally for reconnection with[1–5] and without[6] a guide field. They have also been predicted analytically in electron[7] and Hall[8–10] magnetohydrodynamics (MHD) models *without a guide field*. However, no analytical model for reconnection *with a guide field* is currently available. Since a guide field is normally present and sometimes significantly exceeds the reconnecting in-plane magnetic field, as in magnetic fusion experiments, a nonlinear analytical model of fast reconnection with a guide field is highly desirable.

Extended MHD simulations of guide-field reconnection have demonstrated that fast reconnection regimes[1–5] occur when the sound gyroradius $\rho_s \equiv \sqrt{T_e/m_i}/\Omega_i$ exceeds the diffusion region thickness δ.[1,5,11,12] Here, T_e is the electron temperature, m_i is the ion mass, and $\Omega_i \equiv eB_0/(m_i c)$ is the ion gyrofrequency in the guide field B_0, with c as the speed of light and e as the proton electric charge. The latter conclusion was recently confirmed experimentally.[13] However, despite significant theoretical progress, a number of fundamental questions remains unanswered. In particular, the scaling of the nonlinear reconnection rate with ρ_s appears to be problem dependent: Kleva *et al.*[1] found $E_z \propto \rho_s^0$ for magnetic flux bundle coalescence, Bhattacharjee *et al.*[14] observed $E_z \propto \rho_s$ for tearing modes, Fitzpatrick[3] concluded that $E_z \propto \rho_s^{3/2}$ for the Taylor problem, while Schmidt *et al.*[5] saw $E_z \propto \rho_s^{\alpha}$ for tearing modes and magnetic island coalescence, where α decreased from 1 to 0 as ρ_s increased. Formation of a thin current sublayer inside the current sheet layer, which results in explosive reconnection during early nonlinear stages,[14–18] and its nonlinear saturation is also poorly understood. Simulations show that the total layer thickness remains of order the electron skin depth d_e,[16,19] whereas the sublayer thins with time without limit.[14–18] It has been speculated[16–18] (but not demonstrated) that such thinning is unphysical and will eventually be arrested by three-dimensional nonlinearities or electron viscosity μ_e. Ion viscosity μ is usually believed to play no significant role in the large-guide-field fast reconnection process (see, e.g., Refs. 1, 3, and 5), with the exception of Ref. 20, where μ was found to play a fundamental role in the Petschek resistive MHD reconnection model (which is now recognized to be unphysical[21]).

Here we generalize the analytical treatment proposed in Refs. 7–9 to study the character of fast magnetic reconnection in extended MHD with a *large guide field*. We consider a quasisteady state diffusion region in two dimensions and evaluate its aspect ratio and the corresponding reconnection rate. We include plasma resistivity η, ion viscosity μ, and finite-sound-gyroradius ρ_s effects. Similar to the zero-guide-field case,[7–10] we find that two reconnection regimes are possible: a fast and a slow, Sweet–Parker-like.[22] However, unlike the zero-guide-field case, where the reconnection rate in the fast regime is not explicitly dependent on dissipation coefficients, for large guide fields we find it proportional to $\mathrm{Pr}^{-1/2}$ with $\mathrm{Pr} = \mu/\eta$ the magnetic Prandtl number. Following Ref. 21 in assuming that μ corresponds to the perpendicular ion viscosity (as defined in Ref. 23) gives $\mathrm{Pr}^{-1/2} \sim \beta^{-1/2}(m_e/m_i)^{1/4}(T_i/T_e)^{1/4}$, which predicts the reconnection rate is independent of collision frequencies. Here, $\beta \ll 1$ is the ratio of plasma pressure to magnetic pressure, and m_e and T_i are the electron mass and ion temperature, respectively. A major result of this letter is that the reconnection rate in the fast regime is found to be proportional to $\mu^{-1/2}$, and thus ion viscosity plays a fundamental role. A transition between the two regimes occurs at $\delta \sim \rho_s$, in agreement with previous

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findings.[1,5,12,13] The model gives an explicit dependence of E_z on ρ_s and suggests why different scalings can be observed for different reconnecting systems. It also obtains a nonlinear threshold for the current sheet thinning[14–18] and demonstrates how it is arrested by dissipation processes.

Two-field fluid model. We consider plasma in a Harris magnetic field with a strong, straight, homogeneous guide field $\boldsymbol{B}_0 = B_0\hat{z}$, $\hat{z} \equiv \nabla z$, and employ the two-dimensional $(\partial/\partial z \equiv 0)$, low-$\beta$ $(\beta \ll 1)$, cold-ion $(0 < T_i \ll T_e)$ fluid model[1,5,14,17,21,24] for magnetic reconnection in the x-y plane. The model assumes homogeneous equilibrium plasma density n_0, and homogeneous, time-independent T_i and T_e. Herein, we neglect finite-d_e effects, which will be discussed in future work. The assumption $T_i \ll T_e$ justifies retaining finite-ρ_s effects while neglecting finite-ion-gyroradius effects. The model neglects electron viscosity, which is important in collisionless plasmas and will be included in future work. In addition, a simple ion viscosity closure is used that corresponds to retention of the ion perpendicular viscosity only (see Ref. 21 for a detailed justification). Normalizing the model equations to the Alfvén speed $V_A \equiv B_0/\sqrt{4\pi n_0 m_i}$ and an arbitrary length L gives

$$\left(\frac{\partial}{\partial t} + \boldsymbol{V} \cdot \nabla\right)\varpi = \boldsymbol{B} \cdot \nabla(\nabla^2\psi) + \mu\nabla^2\varpi, \tag{1}$$

$$\frac{\partial \boldsymbol{B}}{\partial t} + \nabla \times [\boldsymbol{B} \times (\boldsymbol{V} - \rho_s^2\hat{z} \times \nabla\varpi)] = -\eta\nabla \times (\nabla \times \boldsymbol{B}). \tag{2}$$

Here, $\boldsymbol{V} \equiv \hat{z} \times \nabla\varphi$ is the ion flow velocity with φ the electrostatic potential, $\varpi \equiv \nabla^2\varphi$ is the flow vorticity, and $\boldsymbol{B} \equiv \hat{z} \times \nabla\psi$ is the magnetic field in the x-y plane. Quantities η and μ are normalized resistivity and ion viscosity, respectively. In numerical simulations, Eq. (2) is often replaced with a single scalar equation for ψ. However, for our purposes, it is convenient to consider the vector form of the equation.[7–9,19]

We concentrate on the diffusion region and do not concern ourselves with a particular physical system supplying magnetic flux. A closed dynamical description can be obtained by coupling our microscopic diffusion region description with a suitable macroscopic driver, as in Ref. 25. As shown in Fig. 1, we assume a rectangular diffusion region of (normalized) dimensions δ and w with plasma entering and exiting along the $\hat{\boldsymbol{y}} \equiv \nabla y$ and $\hat{\boldsymbol{x}} \equiv \nabla x$ directions, respectively. We define the discrete upstream and downstream magnetic field variables $B_x \equiv \hat{\boldsymbol{x}} \cdot \boldsymbol{B}(0, \delta/2)$ and $B_y \equiv \hat{\boldsymbol{y}} \cdot \boldsymbol{B}(w/2, 0)$, respectively, and the discrete flow stream function $\Phi \equiv -\varphi(w/2, \delta/2)$. Then, the inflow and outflow velocities are given by $V_y = -2\Phi/w$ and $V_x = 2\Phi/\delta$, respectively. Following Refs. 7–9 and 19 we discretize Eq. (1) at $(x,y) = (w/2, \delta/2)$ and the $\hat{\boldsymbol{x}}$ and $\hat{\boldsymbol{y}}$ components of Eq. (2) at $(x,y) = (0, \delta/2)$ and $(w/2, 0)$, respectively. Time-derivative terms are normally found small in nonlinearly saturated states, e.g., at and around the time of the local maximum of the reconnection rate.[25] A steady-state analysis of the diffusion region (but not the full domain) is appropriate in such situations. Neglecting time derivatives and numerical factors

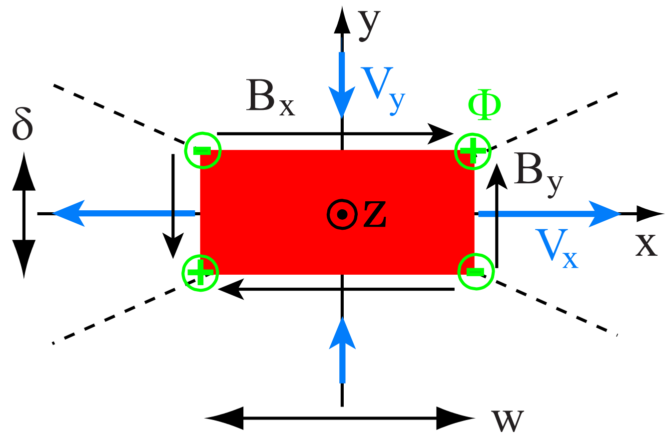

FIG. 1. (Color online) Diffusion region geometry.

of order unity and following a procedure similar to that in Refs. 7–9 and 19 results in the following equations for Φ, B_x, and B_y,

$$\frac{\Phi^2}{\delta w}\left(\frac{1}{\delta^2} - \frac{1}{w^2}\right) + \left(\frac{B_x}{w} + \frac{B_y}{\delta}\right)\left(\frac{B_y}{w} - \frac{B_x}{\delta}\right) = -\mu\Phi\left(\frac{1}{\delta^2} + \frac{1}{w^2}\right)^2, \tag{3}$$

$$-\frac{\Phi B_x}{\delta w}\left[1 + \rho_s^2\left(\frac{1}{\delta^2} + \frac{1}{w^2}\right)\right] = \eta\left(\frac{B_y}{\delta w} - \frac{B_x}{\delta^2}\right), \tag{4}$$

$$\frac{\Phi B_y}{\delta w}\left[1 + \rho_s^2\left(\frac{1}{\delta^2} + \frac{1}{w^2}\right)\right] = \eta\left(\frac{B_x}{\delta w} - \frac{B_y}{w^2}\right). \tag{5}$$

Equations (3)–(5) are invariant under the plasma flow reversal, i.e., under the substitution $(B_x, B_y, \Phi, \delta) \leftrightarrow (B_y, B_x, -\Phi, w)$, as expected.[7,19] They contain five unknowns: δ, w, B_x, B_y, and Φ. It is therefore necessary to consider two independent parameters. It is convenient to choose w and B_x as such parameters.

Solution of the discrete equations. An important characteristic of the diffusion region is its aspect ratio $\xi \equiv \delta/w$. The reconnection rate $E_z = -\eta J_z = \eta(B_x/\delta - B_y/w)$ can be conveniently expressed in terms of ξ as

$$E_{z*} \equiv \frac{E_z}{B_x^2} = \frac{\sqrt{2}(1-\xi^2)}{S_\eta \xi}, \tag{6}$$

where $S_\eta \equiv \sqrt{2}B_x w/\eta$ is the resistive Lundquist number. It is clear from Eq. (6) that for given B_x and w, large reconnection rates preferentially occur for $\xi \ll 1$. We will concentrate next on the $\xi < 1$ limit by approximating $1+\xi^2 \approx 1-\xi^2 \approx 1$ in Eq. (6) and elsewhere. Then, the reconnection rate becomes $E_{z*} \approx \sqrt{2}/(S_\eta \xi)$. An equation for ξ, or equivalently δ, can be obtained from Eqs. (3)–(5). Introducing a dimensionless quantity $\hat{\delta} \equiv \delta/\rho_s$, which characterizes the importance of plasma compressibility, and eliminating variables results in the following equation for ξ:

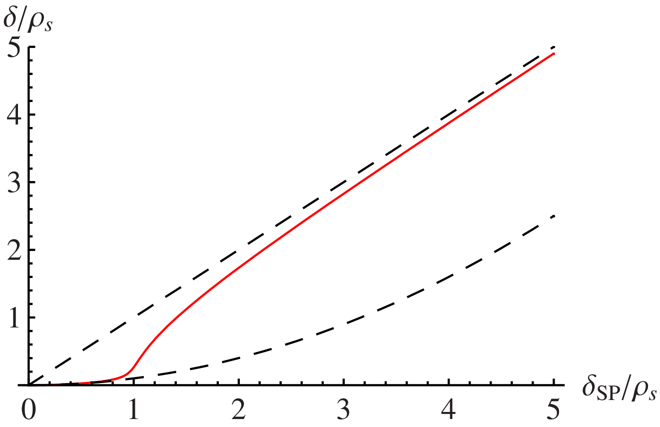

FIG. 2. (Color online) Numerical solution of Eq. (8) with $q=0.01$ (solid line). Asymptotics (9) are shown with dashed lines.

$$\frac{1}{S_\eta^2\xi^4}+\frac{1+\hat{\delta}^2}{S_\eta S_\mu \xi^4\hat{\delta}^2}=\frac{(1+\hat{\delta}^2)^2}{\hat{\delta}^4}, \tag{7}$$

where $S_\mu\equiv\sqrt{2}B_x w/\mu$ is the viscous Lundquist number.

To solve Eq. (7) for ξ, or equivalently $\delta=\xi w$, it is convenient to introduce parameters $q\equiv \mathrm{Pr}/(\mathrm{Pr}+1)\leq 1$, with $\mathrm{Pr}\equiv\mu/\eta$ the magnetic Prandtl number; and $\hat{\delta}_{\rm SP}\equiv\delta_{\rm SP}/\rho_s$, with $\delta_{\rm SP}\equiv(w/\sqrt{S_\eta})(1+\mathrm{Pr})^{1/4}$ as the Sweet–Parker length scale.[20] Then, Eq. (7) becomes

$$\hat{\delta}^2(\hat{\delta}^2+1)^2-\hat{\delta}_{\rm SP}^4(\hat{\delta}^2+q)=0 \tag{8}$$

so that

$$\delta\approx\begin{cases}\delta_{\rm SP} & \delta_{\rm SP}\gg\rho_s\\ w^2/(\rho_s\ \mathrm{Ha}) & \delta_{\rm SP}\ll\rho_s,\end{cases} \tag{9}$$

where $\mathrm{Ha}=\sqrt{S_\eta S_\mu}$ is the Hartmann number. The corresponding reconnection rates are

$$E_{z*}\approx\begin{cases}\sqrt{2/S_\eta}(\mathrm{Pr}+1)^{-1/4} & \delta_{\rm SP}\gg\rho_s\\ \sqrt{2/\mathrm{Pr}}(\rho_s/w) & \delta_{\rm SP}\ll\rho_s.\end{cases} \tag{10}$$

Consequently, two reconnection regimes are possible. When δ exceeds ρ_s, the reconnection is Sweet–Parker-like[22] and therefore slow. In the opposite limit, the diffusion region thickness is proportional to Ha^{-1} and the reconnection rate to $\mathrm{Pr}^{-1/2}=\sqrt{\eta/\mu}$. Therefore, both ion viscosity and resistivity are essential ingredients. However,[21] if μ corresponds to the perpendicular ion viscosity, then $\mathrm{Pr}^{-1/2}\sim\beta^{-1/2}(m_e/m_i)^{1/4}(T_i/T_e)^{1/4}$ and the reconnection rate is not explicitly dependent on collision frequencies. Hence, it is potentially fast. Moreover, since $\rho_s=\sqrt{\beta/2}d_i$ with $d_i=(c/L)\sqrt{m_i/(4\pi n e^2)}$ the normalized ion inertial length scale, the reconnection rate $E_{z*}=(2/\mathrm{Pr})^{1/2}(\rho_s/w)$ in the fast reconnection regime is also independent of $\beta\ll 1$. A transition between the two regimes occurs at $\delta\sim\delta_{\rm SP}\sim\rho_s$, as expected.[1,5,12,13] These conclusions are confirmed by solving Eq. (8) numerically. The solution for $q=0.01$ is shown in Fig. 2 with a solid line. Asymptotics (9) are shown with dashed lines.

Several important conclusions can be drawn from Eqs. (9) and (10). First, a steady-state fast reconnection regime does not exist when $\mu=0$. This confirms the indefinite current sublayer thinning observed in Refs. 14–18 in this regime. Finite ion viscosity is therefore essential to regularize the solution for δ, $\delta_{\rm SP}\lesssim\rho_s$. Second, Eqs. (9) and (10) predict that the fast reconnection regime does not exist when $\rho_s=0$ for arbitrary μ, as expected.

Fundamental role of ion viscosity. The relevance of μ in the fast regime can be understood directly from Eqs. (1) and (2). Consider a quasisteady-state diffusion layer and assume $\xi\ll 1$ so that according to Eqs. (4) and (5), $B_y/B_x=\xi\ll 1$. Equation (1) requires the Lorentz force $\boldsymbol{B}\cdot\boldsymbol{\nabla}(\nabla^2\psi)$ to be balanced by either advection $\boldsymbol{V}\cdot\boldsymbol{\nabla}\varpi$ or viscous diffusion $\mu\nabla^2\varpi$. Therefore, viscosity will become essential when $\mu\nabla^2\sim\mu/\delta^2>\boldsymbol{V}\cdot\boldsymbol{\nabla}\sim\Phi/(\delta w)$, and unimportant otherwise. The magnitude of Φ can be estimated from Eq. (2) by discretizing its x component at $(x,y)=(0,\delta/2)$,

$$\frac{B_x\Phi}{\delta w}\left(1+\frac{\rho_s^2}{\delta^2}\right)\sim\frac{\eta B_x}{\delta^2}. \tag{11}$$

It follows that

$$\frac{\Phi}{\delta w}\sim\begin{cases}\eta/\delta^2 & \delta>\rho_s\\ \eta/\rho_s^2 & \delta<\rho_s.\end{cases} \tag{12}$$

Consequently, when $\delta/\rho_s<\min\{1,\sqrt{\mathrm{Pr}}\}$, viscosity will play a fundamental role in balancing the Lorentz force term at the diffusion region, and thus in the reconnection process. On the other hand, for $\delta>\rho_s$ both terms are in principle comparable and μ does not influence reconnection fundamentally.

Comparison with previous numerical studies. The results obtained herein agree well with the available numerical simulations of guide field reconnection. Specifically, Kleva *et al.*[1] employed Eqs. (1) and (2) with $\mu=\eta$ (i.e., $\mathrm{Pr}=1$) to study coalescence of two magnetic flux bundles. The work varied ρ_s and η in the fast regime ($\delta<\rho_s$) by as much as a factor of four each and found $w\propto\rho_s$ and $E_z\propto\eta^0$, ρ_s^0, consistent with Eq. (10). In addition, it was observed that $\delta\propto\eta$, in agreement with Eq. (9) when $\mu=\eta$.

Schmidt *et al.*[5] also employed Eqs. (1) and (2) with $\mu=\eta$ (again, $\mathrm{Pr}=1$), but studied reconnection driven by a tearing mode and by a magnetic island coalescence instability. In both cases, the reconnection rate in the fast regime ($\delta<\rho_s$) was again found to be independent of η, $E_z\propto\eta^0$. Moreover, it was observed that $E_z\propto\rho_s^\alpha$ with $\alpha\sim 1$ for $\delta\sim\delta_{\rm SP}\sim\rho_s$ (i.e., at the slow-to-fast transition) and $\alpha\sim 0$ for $\rho_s\gg\delta_{\rm SP}$. Both scalings are also consistent with Eq. (10) when one considers that at the transition, w is still determined by the Sweet–Parker dynamics and is therefore independent of ρ_s (i.e., $w\propto\rho_s^0$; hence $E_z\propto\rho_s/w\propto\rho_s$), while $w\propto\rho_s$ when δ is deep in the ρ_s sublayer,[1] resulting in $E_z\propto\rho_s/w\propto\rho_s^0$.

One set of results apparently at odds with our predictions is documented in Ref. 3, where a system of four nonlinear equations for ψ, ϖ, and the $\hat{z}$ components of the perturbed magnetic field and ion flow velocity is employed that is valid for both $\beta\ll 1$ and $\beta\gg 1$ regimes. When solving this system in the nonlinear reconnection regime with a large guide field, it was found that $E_z \propto \eta^0 \rho_s^{3/2}$ with fixed μ, which appears to disagree with our findings (9) and (10). However, no data are provided for w in the reference. Furthermore, these four equations can only be reduced to our Eqs. (1) and (2) when $\beta \ll 1$, $\kappa = \mu \ll \beta\eta$, and $\nu \equiv 0$, with κ the plasma heat conductivity and ν the hyper-resistivity as defined in Ref. 3. Since Ref. 3 employed $\kappa = \mu \sim \beta\eta$ and $\nu = 2.5 \times 10^{-9} > 0$, the results therein cannot yet be compared with our theory.

In conclusion, we have applied the analytical framework developed in Refs. 7–9 to the study of magnetic reconnection with a large guide field. We have confirmed that two reconnection regimes are possible depending on whether the diffusion region thickness δ is larger or smaller than the sound gyroradius ρ_s. For $\delta > \rho_s$ the reconnection is slow and described by the Sweet–Parker expressions, modified to account for ion viscosity μ. For $\delta < \rho_s$, the character of reconnection changes, with δ becoming proportional to Ha^{-1}, with $\mathrm{Ha} = \sqrt{S_\eta S_\mu}$ as the Hartmann number, and the reconnection rate E_z becoming proportional to $\mathrm{Pr}^{-1/2}$, with $\mathrm{Pr} = \mu / \eta$ as the Prandtl number. Assuming that μ describes the perpendicular ion viscosity results in E_z independent of collision frequencies and plasma β, and therefore potentially fast. The transition between the two regimes occurs at $\delta \sim \delta_{\mathrm{SP}} \sim \rho_s$.

One major contribution of this letter is the uncovering of the fundamental role that ion viscosity μ plays in enabling a steady-state, fast-reconnecting, large-ρ_s regime in low-β plasmas. In particular, we have demonstrated that such a regime does not exist when $\mu = 0$. This result explains numerical observations of indefinite current sheet thinning when $\delta < \rho_s$ in the absence of ion viscosity.[14–18]

Given the demonstrated relevance of ion viscosity for the low-β Hall MHD reconnection, it seems clear that the simple resistive and viscous closures commonly used in the literature for Eqs. (1) and (2) are likely to be inadequate for accurately describing weakly collisional magnetospheric and magnetic fusion plasmas. It follows that more accurate models for collisionless electron and ion viscosities (e.g., gyroviscosities[23]) must be employed in Eqs. (2) and (1), respectively, to adequately explain experimental observations. The consideration of such closures in our theoretical framework is left for future work.

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