Magnetic reconnection in partially ionized plasmas

Ellen G. Zweibel,[1,a)] Eric Lawrence,[2,b)] Jongsoo Yoo,[2,c)] Hantao Ji,[2,d)] Masaaki Yamada,[2,e)] and Leonid M. Malyshkin[3,f)]
[1]*Department of Astronomy, University of Wisconsin, 475 North Charter Street, Madison, Wisconsin 53706, USA*
[2]*Princeton Plasma Physics Laboratory, P. O. Box 451, Princeton, New Jersey 08543, USA*
[3]*Department of Astronomy & Astrophysics, University of Chicago, 5640 South Ellis Avenue, Chicago, Illinois 60637, USA*

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We review the theory of magnetic reconnection in weakly ionized gases. The theory is relevant to reconnection in the interstellar medium, protostellar and protoplanetary disks, the outer envelopes of cool stars, and a new laboratory experiment. In general, partial ionization introduces three effects beyond the obvious one: increased resistivity due to electron-neutral collisions. First, magnetic neutral sheets are steepened by plasma-neutral drift, setting up the conditions for reconnection. Second, when ion-neutral friction is strong, the effective ion mass is increased by ρ/ρ_i, the ratio of total to plasma mass density. This reduces the Alfven speed v_A by a factor of $\sqrt{\rho/\rho_i}$ and increases the ion skin depth δ_i by $\sqrt{\rho/\rho_i}$. As a result, entrainment of neutrals slows MHD reconnection but permits the onset of fast collisionless reconnection at a larger Lundquist number S, or for a longer current sheet, than in the fully ionized plasma case. These effects, taken together, promote fast collisionless reconnection when the ionization fraction is of order 10% to 1%, but reconnection is slowed down for much smaller ionization fractions. Finally, ion-neutral friction can be a strong heating mechanism throughout the inflow and outflow regions. These effects are under study at the Magnetic Reconnection Experiment (MRX). © *2011 American Institute of Physics*. [doi:10.1063/1.3656960]

I. INTRODUCTION

Magnetic reconnection is a fundamental process in plasma physics.[1,2] Understanding reconnection in astrophysical plasmas presents serious challenges. The range of relevant lengthscales and timescales is usually extreme, diagnostics are based on remote sensing, and the system may have novel features not encountered in conventional laboratory plasmas.

One such feature is a very low ionization fraction. The ionization fraction in the solar chromosphere and in the so-called warm neutral interstellar medium is of order 10^{-2}, in the dense interstellar clouds where stars are formed it can be as low as 10^{-7}, and in protostellar and protoplanetary accretion disks and the atmospheres of brown dwarf stars it can be 10^{-10} or even less. Although at the lowest ionization fraction the electrical conductivity itself disappears, the large physical size of many astrophysical plasmas makes their Lundquist number S itself very large, and reconnection is the most important process for dissipating magnetic energy and changing the magnetic topology. This point is becoming widely appreciated; see, for example, Refs. 3 and 4. Thus, it is important to understand reconnection in partially ionized gases.

In Sec. II of this paper, we review the theory of reconnection in partially ionized gases, including, in order of development, the effect of ion-neutral friction on MHD tearing modes, the steepening of neutral sheets by ion-neutral friction and fast magnetic merging within them, and conditions for the onset of the Hall effect in steady, neutral sheet reconnection. In Sec. III, we discuss the prospects for studying these processes in the laboratory and show some preliminary results. Section IV is a summary.

II. THEORY

We consider three fluids: electrons, a single species of singly charged ions, and a single species of neutrals. We assume quasineutrality ($n_e = n_i$), and thus need only the continuity equations for ions and neutrals

$$\frac{\partial \rho_i}{\partial t} + \nabla\cdot(\ \rho_i \mathbf{u}_i) = \zeta\rho_n - \alpha\rho_i^2, \tag{1}$$

$$\frac{\partial \rho_n}{\partial t} + \nabla\cdot(\rho_n \mathbf{u}_n) = -\zeta\rho_n + \alpha\rho_i^2, \tag{2}$$

where ζ and α are the ionization and recombination rate coefficients, respectively. The momentum equations for the three species are

$$0 = -n_i e\left(\mathbf{E} + \frac{\mathbf{u}_e \times \mathbf{B}}{c}\right) - \nabla P_e - \rho_e \nu_{en}(\mathbf{u}_e - \mathbf{u}_n) - \rho_e \nu_{ei}(\mathbf{u}_e - \mathbf{u}_i), \tag{3}$$

$$\rho_i \frac{D_i \mathbf{u}_i}{Dt} = n_i e\left(\mathbf{E} + \frac{\mathbf{u}_i \times \mathbf{B}}{c}\right) - \nabla P_i - \rho_i \nu_{in}(\mathbf{u}_i - \mathbf{u}_n) - \rho_i \nu_{ie}(\mathbf{u}_i - \mathbf{u}e), \tag{4}$$

$$\rho_n \frac{D_n \mathbf{u}_n}{Dt} - \nabla P_n - \rho_n \nu_{ni}(\mathbf{u}_n - \mathbf{u}_i) - \rho_n \nu_{ne}(\mathbf{u}_n - \mathbf{u}_e), \tag{5}$$

a)Electronic mail: zweibel@astro.wisc.edu.
b)Electronic mail: elawrenc@pppl.gov.
c)Electronic mail: jyoo@pppl.gov.
d)Electronic mail: hji@pppl.gov.
e)Electronic mail: myamada@pppl.gov.
f)Electronic mail: leonmal@flash.uchicago.edu.

where $D_s/Dt \equiv (\partial_t + \mathbf{u}_s \cdot \nabla)$ and ν_{ab}, the rate at which species a collides with species b, can be written $\nu_{ab} = m_b n_b \langle \sigma v \rangle_{ab}/(m_a + m_b)$. Here, $\langle \sigma v \rangle_{ab}$ is the rate coefficient; σ is the cross section, v the relative velocity in the center of mass frame, and the brackets denote averaging over the velocity distribution. Note that $\rho_a \nu_{ab} = \rho_b \nu_{ba}$, so in a weakly ionized gas $\nu_{ni}/\nu_{in} = \rho_i/\rho_n \ll 1$. In Eqs. (3)–(5), we have represented momentum exchange between species as a simple drag force. More general expressions which account for momentum gained or lost by chemical reactions are given in Ref. 5. We have neglected electron inertia in Eq. (3) and assumed all the pressures are isotropic scalars. It is shown in Ref. 6 that the drag of electrons on neutrals can be ignored in calculating the bulk momentum, although electron-neutral collisions affect the resistivity.

Adding Eqs. (3) and (4), neglecting electron-neutral drag compared with ion-neutral drag,[6] and using $\mathbf{J} \equiv en_i(\mathbf{u}_i - \mathbf{u}_e)$ gives the plasma momentum equation

$$\rho_i \left(\frac{\partial}{\partial t} + \mathbf{u}_i \cdot \nabla \right) \mathbf{u}_i = -\nabla(P_e + P_i) + \frac{\mathbf{J} \times \mathbf{B}}{c} - \rho_i \nu_{in} (\mathbf{u}_i - \mathbf{u}_n), \tag{6}$$

used in all MHD treatments.

Equations (1)–(6) have several especially important limits for the reconnection problem. The effect of partial ionization on the MHD tearing mode in slab geometry, with a strong guide field, was first calculated in Ref. 7. In that case the motions are incompressible and 2D. Assuming an $e^{\gamma t}$ time dependence for the tearing perturbations, it can be shown that the effect of ion-neutral collisions on the plasma dynamics is to multiply the plasma inertia term by a factor of $1 + \nu_{in}/(\gamma + \nu_{ni})$. This leads to the tearing mode dispersion relation

$$\gamma^5 \left(1 + \frac{\nu_{in}}{\gamma + \nu_{ni}} \right) = \gamma_*^5, \tag{7}$$

where γ_* is the tearing mode growth rate in the plasma alone (but with electron-neutral collisions included in the resistivity). Equation (7) shows that the stability criterion for tearing modes is unaffected by ion-neutral friction, but the growth rate is affected. In a weakly ionized system, if the growth time γ^{-1} is longer than the neutral-ion collision time ν_{ni}^{-1}, the plasma and the neutrals are well coupled, and $\gamma \approx \gamma_* (\rho_i/\rho_n)^{1/5}$. This result could be derived by replacing the plasma Alfven speed $v_{Ai} \equiv B/\sqrt{4\pi\rho_i}$ by the bulk Alfven speed $v_A \equiv B/\sqrt{4\pi\rho}$, assuming $\rho \approx \rho_n$. On the other hand, if $\gamma \gg \nu_{in}$, the growth time is shorter than the ion-neutral collision time ν_{in}^{-1}, and neutrals have no effect on the tearing mode. Similar estimates for steady reconnection are also presented in Ref. 7.

Ion-neutral collisions have a much larger effect on reconnection in neutral sheets, with no guide field present. If the ionization fraction is very low, and collisions are strong, it is permissible to neglect plasma pressure and inertia in Eq. (6) and simply balance the Lorentz force with ion-neutral friction. This treatment, which is analogous to balancing the Lorentz and drag forces on electrons to derive Ohm's law in resistive MHD, was originally proposed by Mestel and Spitzer.[8] The resulting drift velocity is

$$\mathbf{u}_D \equiv \mathbf{u}_i - \mathbf{u}_n = \frac{\mathbf{J} \times \mathbf{B}}{\rho_i \nu_{in} c}. \tag{8}$$

Again, neglecting plasma inertia, the center of mass velocity $\mathbf{u}$ is approximately the neutral velocity $\mathbf{u}_n$. Using Eq. (8), one can then write the ion velocity as $\mathbf{u}_i \approx \mathbf{u}_D + \mathbf{u}$. When the Hall effect is unimportant (it can easily be included), the magnetic induction equation can then be written as

$$\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{u} \times \mathbf{B}) + \nabla \times \left(\frac{\mathbf{J} \times \mathbf{B}}{\rho_i \nu_{in} c} \times \mathbf{B} \right) + \eta \nabla^2 \mathbf{B}. \tag{9}$$

The first term on the right hand side of Eq. (9) represents induction by the bulk flow $\mathbf{u}$. It is fictitious in the sense that it is the plasma flow, not the bulk, primarily neutral fluid flow that is responsible for induction. The second term represents the difference between the bulk flow and the plasma flow and corrects the first term. Because the plasma and neutral flows would be the same in the absence of magnetic forces, the second term is nonlinear in B. The third term is Ohmic diffusion (simplified by the assumption of constant resistivity η). The drift of magnetic field with respect to the bulk flow, represented by the second term, was dubbed "ambipolar diffusion" by Mestel and Spitzer.[8]

Equation (9) can also be written in terms of the Cowling resistivity (Ref. 9; see also Refs. 10 and 11)

$$\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{u} \times \mathbf{B}) - c\nabla \times (\eta \mathbf{J}_\parallel + \eta_c \mathbf{J}_\perp), \tag{10}$$

where $\mathbf{J}_{\parallel,\perp}$ are the projections of $\mathbf{J}$ parallel and perpendicular to $\mathbf{B}$, respectively, and $\eta_c \equiv \eta + B^2/(\rho_i \nu_{in} c^2)$ is the Cowling resistivity in the limit $\rho_i/\rho_n \ll 1$. To allow for arbitrary neutral fraction, $\eta_c - \eta$ should be multiplied by a factor of $(\rho_n/\rho)^2$.

Equation (9) takes a particularly simple form when $\mathbf{u} = 0$ and the fieldlines are straight, say $\mathbf{B} = \hat{z}B(x,t)$. Then, Eq. (9) becomes

$$\frac{\partial B}{\partial t} = \frac{\partial}{\partial x} \frac{B^2}{4\pi\rho_i \nu_{in}} \frac{\partial B}{\partial x} + \eta \frac{\partial^2 B}{\partial x^2}. \tag{11}$$

Equation (11) shows that in slab geometry, ambipolar drift acts like nonlinear diffusion, with diffusivity v_{Ai}^2/ν_{in}. If B has a null plane, the magnetic field gradient near it becomes large. In the absence of resistivity, Eq. (11) has a steady state solution of the form $B \propto x^{1/3}$. It was shown in Ref. 12 that an initially linear magnetic profile $B \propto x$ steepens to $B \propto x^{1/3}$, and that the component reversals of a magnetic field passively sheared by differential rotation are steepened in the same way. The physics of the steepening is straightforward: the magnetic pressure gradient drives the plasma toward the null; the fieldlines are carried along, and the magnetic profile steepens further.

The $B \propto x^{1/3}$ solution has singular current and infinite drift $\mathbf{u}_D$ at $x = 0$. The effects of resistivity and the buildup of a plasma pressure gradient remove this singularity. However, if these effects are small, the electric current will still become large, and the fieldlines can merge rapidly. This situation was considered in Ref. 13, where it was shown that under conditions typical in the interstellar medium, the ion pressure buildup is limited by rapid recombination, leaving resistivity as the primary effect which removes the singularity. In this parameter regime, the magnetic merging rate is essentially the Sweet-Parker reconnection velocity $(\eta v_{Ai}/L)^{1/2}$, but with the outflow time L/v_{Ai} replaced by the recombination time; a similar result was found in Ref. 14. Under these conditions, magnetic merging can be a very fast process. However, because recombination removes only plasma pressure, not magnetic pressure, even a weak guide field quenches rapid merging.[15]

Collisionless, or Hall mediated, reconnection has received much attention as a possible mechanism for fast reconnection. Experiment, theory, and simulation for a fully ionized plasma have all shown that when the Sweet-Parker reconnection layer thickness $\delta_{SP} = L/\sqrt{S} = \sqrt{L\eta/v_{Ai}}$ is less than the ion skin depth δ_i, the Hall effect comes into play[1,2] (here, L is the length of the current sheet and we have written v_{Ai} to emphasize that this condition has been studied primarily for a fully ionized plasma). Writing the diffusivity η in terms of the electron skin depth δ_e and electron collision time with all species τ_e as $\eta = \delta_e^2/\tau_e$ and using $v_{Ai} = \omega_{ci}\delta_i$, This criterion can be written as

$$L < \delta_i \omega_{ce} \tau_e, \tag{12}$$

for collisionless reconnection, where ω_{ce} is the electron gyrofrequency and collision time, respectively. It is shown in Ref. 16 that Eq. (12) can also be written in terms of the electron mean free path λ_{mfp} as

$$L < 0.48\lambda_{mfp}\left(\frac{m_i}{m_e}\right)^{1/2}, \tag{13}$$

where λ_{mfp} is estimated assuming the plasma pressure in the reconnection layer is equal to the upstream magnetic pressure. Equation (13) makes the connection between Hall reconnection and collisionality more explicit.

Recently, the steady state reconnection equations, allowing for ion-neutral friction and the Hall effect, have been analyzed[6] according to the method developed in Ref. 17. This allows a more precise analysis of steady state reconnection than in Ref. 7 and can be generalized to include the Hall effect. In this work, it was shown that the dynamical effects of ion-neutral friction can be represented by a normalized effective ion mass $\tilde{m} \equiv m_{eff}/m_i$. When the ion-neutral collision time ν_{in}^{-1} exceeds the reconnection outflow time at the ion Alfven speed L/v_{Ai}, $\tilde{m} \sim 1$, and collisions with neutrals have a little effect on MHD reconnection or on the criterion for the onset of Hall reconnection. When the neutral-ion collision time ν_{in}^{-1} is less than the reconnection outflow time at the bulk Alfven speed L/v_A, $\tilde{m} \approx \rho_n/\rho_i$, the effective Alfven speed is $v_A = v_{Ai}/\sqrt{\tilde{m}}$, and the effective ion skin depth becomes $d_i\sqrt{\tilde{m}}$. Equations (12) and (13) generalize to

$$L < \delta_i \omega_{ce} \tau_e \sqrt{\tilde{m}}, \tag{14}$$

and

$$L < 0.48\lambda_{mfp}\left(\frac{m_i \tilde{m}}{m_e}\right)^{1/2}, \tag{15}$$

for Hall mediated reconnection. The result that ion-neutral friction increases the critical current sheet length for onset of Hall mediated reconnection could be important for astrophysics, because the critical current sheet lengths are otherwise quite short.

The intermediate regime $v_{Ai}/\nu_{in} < L < v_{Ai}/\nu_{in}\sqrt{\rho/\rho_i}$ smoothly connects the cases of weak ($L < v_{Ai}/\nu_{in}$) and strong ($L > v_A/\nu_{ni}$) ion-neutral coupling. The rate of frictional heating $H_{in} = \rho_i \nu_{in} v_D^2$ is very large in the intermediate regime and accounts for a substantial fraction of the energy dissipated in reconnection.

Thus, we have shown that ion-neutral friction introduces a second characteristic Alfven speed, based on the total density, into the reconnection problem. This, by itself, slows reconnection down, and is the primary effect on the MHD tearing mode. Neutral sheet reconnection, however, can be speeded up, because the current layer is steepened. And finally, through increasing the ion skin depth, collisions promote the onset of fast reconnection—although the effect of a guide field has not yet been considered. Which of these effects can be observed in the laboratory?

III. LABORATORY EXPERIMENT

An experimental campaign[18] to study reconnection in partially ionized plasmas is underway at the Magnetic Reconnection Experiment (MRX) facility.[19] One focus of our studies is to evaluate the effect of ion-neutral coupling on reconnection. MRX already has a number of diagnostics for measuring the magnetic field and plasma density and temperature, but we have recently performed spectroscopic measurements of the neutral density. We measure the passive emission from neutral lines (typically H_β and D_β or various He lines) using a optical probe previously developed for ion spectroscopy studies.[20] The probe has a line of sight limited to a length of 2–5δ, where δ is the current sheet half-width, allowing us to make localized measurements around the current sheet.

The spectrometer has been calibrated to read absolute spectral radiance, and so we can compute the integrated emissivity of the spectral line and the density of excited atoms in the collection volume. Population ratio coefficients calculated using a collisional-radiative model[21] are then used to determine the density of ground state atoms (i.e., the neutral density). Additionally, neutral temperature and flow can be measured from standard Doppler broadening and line shifts.

In Table I, we show shot-averaged plasma and neutral parameters for three initial hydrogen fill pressures. Fill pressure can be controlled to within 0.2 mTorr, allowing us to closely scan n_n/n_e from 3 to35. Note that the neutral density is not a linear function of the static fill pressure. Processes such as neutral pumping[22] likely reduce the neutral density in the core plasma. Spatial profiles of neutral pressure and

TABLE I. Typical plasma and neutral parameters for high gas fill pressure MRX discharges.

P (mTorr)	21	25	29
n_n (cm^{-3})	1.1×10^{14}	1.6×10^{14}	2.5×10^{14}
n_e (cm^{-3})	3.3×10^{13}	3.6×10^{13}	0.72×10^{13}
T_e (eV)	6.8	7.9	7.8
v_A (cm/s)	6.3×10^{6}	6.4×10^{6}	5.0×10^{6}

their role in the reconnection process will be investigated in future experiments.

Characterizing the effect of neutrals on quantities such as the reconnection rate, effective resistivity, and ion-electron scale separation is still in progress, but we can use the neutral density measurements to determine accessible parameter regimes. As described in Sec. II, ion-neutral coupling is determined by the relative value of the current sheet length L to the two parameters given in Sec. II: v_{Ai}/ν_{in} and $v_{An}/\nu_{ni} = (\rho_n/\sqrt{\rho_i(\rho_i+\rho_n)})v_{Ai}/\nu_{in}$.

To estimate ν_{in}, we take the cross section to be $\sigma_{in} = 7-8 \times 10^{-15}$ cm^2 and independent of energy. This is based on tabulated values of cross sections for charge exchange[23] and momentum transfer.[24] In this case, the Maxwellian-averaged collision frequency is given by

$$\nu_{in} = \frac{m_{in}}{m_i} n_n \sigma_{in} \sqrt{\frac{8T}{\pi m_{in}}}, \tag{16}$$

where m_{in} is the reduced mass $m_i m_n/(m_i + m_n)$, and T is the ion and neutral temperature ($\sim$3 eV in these cases).

In Table II, we show the relative length scales for a two sets of experimental parameters. The current sheet length L was determined using 2D magnetic probe array measurements. In both cases, we are at least in the intermediate coupling regime and should be able to produce conditions in the strong coupling regime.

Increasing the neutral density has two competing effects on the Hall scale length in Eq. (14): increased electron-neutral collisionality will decrease the Hall scale length, while the larger effective ion mass will increase it. Plasma parameters then determine which effect is dominant. The net effect can be seen in the ratio of Eqs. (12) and (14):

$$\frac{\tau_e}{\tau_{ei}} \sqrt{\frac{\rho}{\rho_i}}, \tag{17}$$

where τ_{ei} is the electron collision time with ions only and we have assumed strong coupling ($\tilde{m} \sim \rho/\rho_i$). We plot the values of this ratio for MRX parameters as a function of n_n/n_e in

TABLE II. Measured current sheet length L and ion-neutral coupling parameters for two sample experimental conditions. In the first case ions and neutrals are expected to be strongly coupled ($v_{An}/\nu_{ni} < L$), while the coupling is intermediate ($v_A/\nu_{in} < L < v_{An}/\nu_{ni}$) in the second case.

P (mTorr)	v_A/ν_{in} (cm)	v_{An}/ν_{ni} (cm)	L (cm)
21	4.0	6.3	8
29	1.4	8.1	5

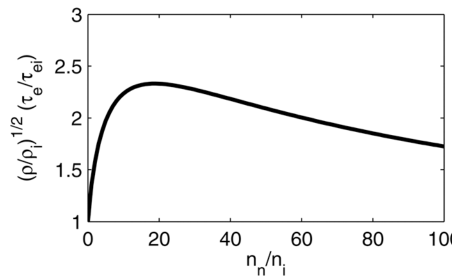

FIG. 1. Hall parameter ratio [Eq. (17)] vs. n_n/n_i for $T_e = 8$ eV, $n_i = 3 \times 10^{13}$ cm^{-3}.

Figure 1. Collision times were determined using classical values for electron-ion collisions[10] and tabulated cross sections for elastic electron-hydrogen collisions.[25] Based on this analysis, the addition of neutrals to MRX plasmas should increase the critical current sheet length by as much as a factor of $\sim$2 at $n_n/n_i \sim$20.

It should be noted that for the plasma parameters given above, we already have $\delta_{SP} < \delta_i$ without the addition of neutrals. Therefore, we plan to examine how different parameters scale as the effective mass is increased well beyond the fully ionized case. Refinement of the plasma production techniques may allow us to create conditions with $\delta_i < \delta_{SP} < \delta_i\sqrt{\tilde{m}}$ in the future.

In Ref. 6, an expression was given for the normalized rate of heating due to ion-neutral friction

$$\frac{Q_{in}}{E_m} \approx \frac{\tilde{\nu}(1-\tilde{v}^2)}{1+\tilde{\rho}\tilde{v}^2}, \tag{18}$$

where Q_{in} is ion-neutral frictional heating, E_m is the rate of magnetic energy supplied to the reconnection layer, $\tilde{\rho} = \rho_n/\rho_i$, $\tilde{\nu}$ is ν_{in} normalized to the ion outflow velocity gradient v_i/L_i, and $\tilde{v}$ is the ratio of the neutral to ion outflow velocity gradients. In the weakly coupled case, the right hand side of Eq. (18) is small because $\tilde{\nu}$ is small, while in the strongly coupled case, it is small because $\tilde{v} \sim 1$. In the intermediate coupling case, the heating rate is predicted to be relatively large since there is a significant frictional slip between the ions and neutrals. We find that for the intermediate coupling case shown above ($P = 29$ mTorr), this equation predicts that $\sim$10% of the supplied magnetic energy will be dissipated through ion-neutral friction. This may be experimentally detectable, and we plan to study such heating in the future.

IV. CONCLUSIONS

Partial ionization is a common state in astrophysical plasmas. For many purposes, these multi-species gases can be treated as a single fluid. However, because magnetic reconnection is essentially a boundary layer phenomenon, the plasma and the neutrals can decouple on the reconnection scale. If the ionization fraction is low, partial decoupling can lead to the formation of intense current layers around magnetic neutral sheets, which facilitates rapid magnetic merging. Other effects—the effect of partial ionization of tearing mode growth rates, the rate of steady MHD reconnection, and the criterion for the transition from MHD to Hall-mediated reconnection—can be understood in terms of an increase in the effective inertia, or mass, of the ions. If the reconnection outflow time is shorter than the ion-neutral collision time ν_{in}^{-1}, neutrals have little effect, and m_{eff} is simply m_i. If the reconnection time is longer than the neutral-ion collision time ν_{in}^{-1}, coupling is strong and m_{eff} is $m_i\rho/\rho_i$. In the intermediate coupling regime, the m_{eff} increases smoothly between these two limits, and frictional heating is strong.

Because the ion skin depth δ_i scales as $m_{eff}^{1/2}$ while the width of the Sweet-Parker layer δ_{sp} scales as $m_{eff}^{1/4}$, the usual criterion for the onset of collisionless reconnection, $\delta_i/\delta_{SP} > 1$, is modified such that maximum current sheet length at which collisionless reconnection can occur scales as $m_{eff}^{1/4}$. Although this enlargement of the parameter space for collisionless reconnection is to some degree offset by increased resistivity due to electron-neutral collisions, it may be possible to measure this effect in an experiment currently under way at MRX.

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