Reconnection in the ISM

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ABSTRACT

We discuss the role of ambipolar diffusion for simple reconnection in a partially ionized gas, following the reconnection geometry of Parker and Sweet. When the recombination time is short the mobility and reconnection of the magnetic field is substantially enhanced as matter escapes from the reconnection region via ambipolar diffusion. Our analysis shows that in the interstellar medium it is the recombination rate that usually limits the rate of reconnection. Consequently, the typical reconnection velocity in interstellar medium is \(\sim (\eta/\tau_{\text{recomb}})^{1/2}(1 + x^{-1} \beta^{-1})\), where \(\eta\) is the ohmic resistivity, \(x^{-1}\) is the ionization fraction, and \(\beta\) is the ratio of gas pressure to magnetic pressure. We show that heating effects can reduce this speed by increasing the recombination time and raising the local ion pressure. In the colder parts of the ISM, when temperatures are \(\sim 10^2\) K or less, we obtain a significant enhancement over the usual Sweet-Parker rate, but only in dense molecular clouds will the reconnection velocity exceed \(10^{-3}\) times the Alfvén speed. The ratio of the ion orbital radius to the reconnection layer thickness is typically a few percent, except in dense molecular clouds where it can approach unity.

We briefly discuss prospects for obtaining much faster reconnection speeds in astrophysical plasmas.

Subject headings: Magnetic fields; Galaxies: magnetic fields, ISM: molecular clouds, magnetic fields

1. Introduction

Understanding the mobility and reconnection of magnetic fields in a conducting medium is critical to understanding the evolution and origin of large scale magnetic fields in astrophysical objects such as stars, galaxies, and accretion disks. Standard dynamo theory relies on the turbulent transport of magnetic flux to move the field lines and, implicitly,
to change the large scale topology of the magnetic field. However, the substitution of turbulent transport for microscopic diffusion is difficult to justify on theoretical grounds. Normally we would expect flux freezing to be an excellent approximation to the motion of the magnetic field in a highly conducting fluid. Getting field lines to pass through one another, or rearrange their connections, must ultimately involve ohmic diffusion. Moreover, as the field lines are stretched, they will strengthen and exert forces that will prevent further stretching or deformation. Recently several different authors (e.g. Cattaneo and Vainshtein 1991, Gruzinov and Diamond 1994) have argued that this raises an insuperable obstacle to turbulent transport, and standard mean-field dynamo theory, in highly conducting fluids. One of us (Vishniac 1995) has proposed that this problem may be solved through the formation of intense flux tubes, but this assumes that reconnection is rapid enough to allow the formation of such structures in a small number of dynamical time scales and that the plasma has a high $\beta$ so that the magnetic field can be distributed intermittently.

Estimates of reconnection speeds based on a simplified geometry (Sweet 1958, Parker 1957) give

$$V_{\text{rec}} \approx \left( \frac{v_A \eta}{L_x} \right)^{1/2} = v_A Re_B^{-1/2}$$

where $L_x$ is the typical structure scale, $v_A$ is the Alfvén speed, $\eta$ is the ohmic diffusion constant, and $Re_B \equiv (v_A L_x / \eta)$ is the magnetic Reynolds number. In general, we expect that $v_A$ is comparable to the local turbulent velocity, so that this speed will be many orders of magnetic slower than typical fluid velocities in astrophysical objects, where $Re_B >> 1$. However, we note that observations of magnetic fields in the solar chromosphere and corona (cf. Dere 1996, Innes et al. 1997) suggest that reconnection can occur at speeds $\sim 0.1 v_A$. It is not clear what conditions are necessary for such rapid reconnection or what mechanism is operating that allows it.

In some ways, these problems are especially severe when we consider the evolution of the galactic magnetic field (Kulsrud & Anderson 1992). The huge scales involved limit the number of dynamical time scales available for generating a large scale field and strong theoretical objections have been raised against the possibility of a primordial field strong enough to eliminate the necessity for a galactic dynamo. In addition, the magnetic Reynolds number in the galactic disk is of order $10^{20}$, so that the Sweet-Parker reconnection speed is negligibly small.

On the other hand, the interstellar medium is only partially ionized, and in dense, cool clouds the ionization fraction is very small. This raises the possibility that magnetic field lines which move with the ion velocity may nevertheless be capable of a high degree of mobility relative to the bulk of the gas. Ultimately, the speed of magnetic reconnection is limited by the width of the current sheet dividing magnetic fields of sharply different
orientation. Although importance of ambipolar diffusion has been long stressed (see Mestel 1985), the results obtained so far are contradictory. Recent one dimensional numerical work by Brandenburg and Zweibel (1995) and Zweibel and Brandenburg (1997) suggests that ambipolar diffusion can give rise to narrow current sheets, whose widths are ultimately determined by the size of particle orbits in the plasma. This in turn suggests that reconnection in the neutral parts of the interstellar medium may play a critical role in the galactic dynamo (cf. Subramanian 1998). At the same time analytical studies of reconnection in the presence of ambipolar diffusion by Dorman (1996), also based on a one dimensional model, led to a different conclusion.

These ambiguities motivate our current study. In this paper we reexamine the role of ambipolar diffusion in Sweet-Parker reconnection, including the effects of ion pressure and finite recombination rates for the ions, and allowing for the transverse loss of neutral particles during reconnection. We calculate reconnection velocities and show that for most astrophysically important cases reconnection is limited by the recombination rate. We find a substantial enhancement of the reconnection rates even in media with low rates of recombination. This enhancement is not enough, however, to account for magnetic flux tube formation in the ISM (as discussed in Lazarian & Vishniac (1996) and Subramanian (1998)). At the same time the rates obtained here can explain topological changes that accompany stellar formation, that is, the disconnecting of the magnetic field of a collapsing cloud from the interstellar magnetic field.

In section II of this paper we will briefly review the physical basis for the Sweet-Parker rate, and for the enhancement due to ambipolar diffusion. In section III we estimate the effects of ion pressure and give a simple mathematical model supporting our claims. In section IV we calculate rates of reconnection for different phases of the interstellar medium. Finally, in section V we discuss the implications and limitations of our work, and summarize our conclusions.

2. Sweet-Parker Reconnection and Ambipolar Diffusion

2.1. The Case of the Imperfect, Ionized Gas

We start by reviewing the physical basis for the Sweet-Parker reconnection rate. This is well known material, but including it here will make subsequent arguments concerning ambipolar diffusion and recombination clearer. The simple geometry that forms the basis of the Sweet-Parker reconnection rate consists of two regions of opposite magnetic polarity facing one another. The $\hat{y}$ axis is perpendicular to the field lines and gas is approaches the
midplane with a velocity $-u_y$ for $y > 0$ and $u_y$ for $y < 0$. The magnetic field lines curve apart a distance $\pm L_x$ along the $\hat{x}$ axis. As the field lines curve apart the gas streams to the right and left with a velocity $u_x$. The magnetic field is zero along the $\hat{x}$ axis, but rises to $\pm B$ a distance $L_y$ on either side. This geometry is illustrated in figure 1, although the arrows for gas flow are appropriate for the case of ambipolar diffusion. When the gas is completely ionized the flow along the $\hat{x}$ axis is confined to the layer of width $\Delta$.

The evident neglect of the existence of a third dimension is the major weakness of this picture. However, the tension of the magnetic field lines, and their consequent resistance to bending, makes the role of this neglected dimension quite complicated. Here we will simply assume that this geometry provides a useful approach to the problem of reconnection, and defer consideration of more complicated geometries to a subsequent paper. It is important to keep in mind that in three dimensions the magnetic field along the reconnection surface does not actually vanish, but simply reduces to a component in the $\hat{z}$ direction which is common to both magnetized regions. This does not effect our arguments in any way, but implies that one should resist the temptation to think of the plasma in the reconnection zone as essentially unmagnetized.

Conservation arguments can be used to estimate $u_x$ and constrain the reconnection geometry. First, we note that the pressure along a line parallel to the $\hat{y}$ axis and passing through the origin is approximately constant, i.e.

$$P + \frac{B^2}{8\pi} \approx \text{constant}. \quad (2)$$

At the midplane the magnetic pressure vanishes. Following the fluid to the right or left, one comes (in a distance $\sim L_x$) to a region unaffected by the magnetic pressure. This implies a pressure excess in the reconnection region of $B^2/8\pi$. This pressure excess is sometimes derived by considering the dissipation of magnetic energy in the reconnection layer, but this gives rise to the mistaken notion that an efficiently radiating reconnection layer will not have a gas pressure excess. From Bernoulli’s theorem

$$u_x^2 + \frac{P}{\rho} \approx \text{constant}, \quad (3)$$

which implies that if $u_x = 0$ at the origin then at $x = \pm L_x$

$$u_x = \left(\frac{\Delta P}{\rho}\right)^{1/2} \approx \frac{B}{\sqrt{8\pi\rho}} = \frac{v_A}{2^{1/2}}, \quad (4)$$

where $v_A$ is the Alfvén velocity of the magnetized regions.
The mass influx into the reconnection region is just \(2\rho L_x u_y\), while the mass efflux is \(2\rho L_x u_y = 2\rho L_y u_x\), so from conservation of mass we have

\[
u_y = \frac{L_y}{L_x} v_A.
\] (5)

We note that the ions need not trace the movement of the magnetic flux in an imperfect fluid. The magnetic field current is

\[
J_B = v_i \otimes B - \eta \nabla B,
\] (6)

where the subscript \(i\) refers to the ions. Given the geometry used here, the effective velocity of the magnetic field can deviate from \(u_y\) by an amount \(\eta/L_y\). To put it another way, the magnetic field is no longer flux frozen if \(L_y\) is sufficiently small and regions of opposite polarity can annihilate much faster than matter can be dragged into the reconnection region. This acts as a kind of regulator for the size of the reconnection region. If it is too large, then reconnection slows to a crawl and whatever external forces are pushing the magnetized regions together will continue to do so\(^1\). If it is too small, then reconnection runs away and broadens the neutral zone. We can therefore take \(\eta \approx L_y u_y\) and obtain

\[
u_y \sim \left(\frac{\eta v_A}{L_x}\right)^{1/2}.
\] (7)

This is the Sweet-Parker reconnection rate given in equation (1).

One important aspect of this model is the existence of a pressure deficit, of order the magnetic pressure, at a distance \(\pm L_x\) along the \(\hat{x}\) axis. In a turbulent medium with equipartition between the magnetic and turbulent energies such pressure excesses will come and go on a time scale not much longer than \(L_x/v_A\), so that this stationary model can be no more than a very approximate guide to the structure of the reconnection layer. However, the existence of magnetic tension guarantees that if the magnetic field reaches a locally persistent equilibrium, then we can naturally expect such pressure fluctuations to last as long as the magnetic structure itself. In other words, the tension in magnetic ‘knots’ is sufficient to guarantee a downstream pressure deficit of order \(B^2/8\pi\). One might suppose that the accumulation of ejected ions will erase this pressure deficit, but as reconnection proceeds the reconnected magnetic field lines will be pulled away from the reconnection region, bearing with them the ejected ions. It is also true, but not obvious, that ejected neutrals can usually be ignored in this model, a point we will return to later.

\(^1\)The role of the external forcing is, in fact, more complex. It can squeeze the material from the reconnection zone and thus increase the local Alfvén velocity.
2.2. The Case of the Imperfect, Partially Ionized Gas

When the fluid is partly neutral the charged and neutral particles will no longer move together, and \( \mathbf{v}_i \) is no longer the same as the bulk velocity of the fluid. We can estimate \( \mathbf{v}_i - \mathbf{v}_n \) by balancing the pressure exerted on the ions with the collisional drag due to collisions with the neutrals. We get

\[
\frac{\nabla B^2}{8\pi} + c_i^2 \nabla \rho_i = \rho_i \frac{\mathbf{v}_n - \mathbf{v}_i}{t_{i,n}} = \rho_n \frac{\mathbf{v}_n - \mathbf{v}_i}{t_{n,i}} = -c_n^2 \nabla \rho_n ,
\]

where \( t_{i,n} \) and \( t_{n,i} \) are the collision times for ions impacting on neutrals and vice versa, respectively, and \( c_i \) and \( c_n \) are the ion and neutral sound speeds. If we neglect the ion pressure gradient (more on this later) we can write the neutral particle velocity as

\[
\mathbf{v}_n = \mathbf{v}_i + \frac{\nabla B^2}{8\pi \rho_n} t_{n,i} ,
\]

or

\[
\mathbf{v}_i = \mathbf{v} - \frac{\nabla B^2}{8\pi \rho} t_{n,i} ,
\]

where \( \rho \) and \( \mathbf{v} \) are the total density and bulk velocity of the gas. This implicitly assumes that the neutrals are moving relatively slowly, and dominate the fluid mix, so that the neutral pressure gradient can balance the magnetic pressure when the ion pressure gradient fails to do so. We can use this expression in equation (8) to write the magnetic field current as

\[
\mathbf{J}_B = \mathbf{v} \otimes \mathbf{B} - (\eta + v_A^2 t_{n,i}) \nabla \mathbf{B} ,
\]

where \( v_A \) is the Alfvén velocity relative to the total gas density. From here on we will define \( \eta_{ambi} \equiv v_A^2 t_{n,i} \). Typically \( v_A \sim 10^5 \text{ cm/sec} \). Since \( t_{n,i} \sim 4.8 \times 10^8 / n_i \text{ seconds at low temperatures (Dalgarno 1958)} \), and \( \eta \) is of order \( 10^9 \text{ cm}^2/\text{sec} \), \( \eta_{ambi} \) will be many orders of magnitude larger than \( \eta \) in cool, low density regions in the interstellar medium.

Equation (11) seems to imply that we can use \( \eta_{ambi} \gg \eta \) in place of the usual expression for Ohmic diffusion. The neutral gas will stream outward in the \( \hat{x} \) direction with a velocity \( \sim v_A \), driven by the pressure gradient created by collisions with the in-flowing ions. In reality this derivation includes assumptions which impose severe constraints on the substitution of \( \eta_{ambi} \) for \( \eta \). The most important is that we have assumed that the ion pressure gradient is negligible compared to the magnetic pressure gradient. However, if the ions move inward at a steadily increasing speed towards the reconnection surface then for \( y \ll L_y \) the only escape route for the accumulating particles is to join the opposing flow of neutrals. This can only be true if the recombination time is short, so that ions and neutrals change identity easily. This expression is commonly used in calculating the ambipolar
diffusion time for molecular clouds, in which case one only requires the recombination time to be no greater than the time it takes the magnetic field to drift out of a cloud.

Here we are concerned with the formation of a reconnection layer, in which the relevant time scales are much shorter. We begin by assuming rapid recombination. In that case we can see that the width of the reconnection region, $L_y$, is set by the matter diffusion, as above, by equating the drift velocity to $u_y$. This implies

$$u_y \sim \left( \frac{v_A^3 \eta_{n,i}}{L_x} \right)^{1/2} = \left( \frac{v_A \eta_{ambi}}{L_x} \right)^{1/2},$$

when $\eta_{ambi} > \eta$. Later on we will refer to this expression for $u_y$ as $V_{\text{max}}$ to emphasize that this is the maximal velocity of reconnection that is obtainable through ambipolar diffusion. By the same reasoning we have

$$L_y \sim \left( \frac{\eta_{ambi} L_x}{v_A} \right)^{1/2}. \quad (13)$$

In obtaining expressions (12) and (13) we assumed, first, that the outflow from the region $L_y$ happens with a velocity $v_A$ and, second, that on leaving the reconnection layer neutrals diffuse slowly out while moving together with ions over the distance comparable with $L_x$. The former assumption depends on the pressure differential around the reconnection region, which is assumed as part of our basic geometry, and the dynamics of the outflow, which we check below by studying the structure of the ambipolar diffusion layer. The latter assumption constrains the geometries of the reconnection layers considered.

It’s important to note that the actual process of reconnection takes place only in a narrow zone where $\eta > \eta_{ambi}$. However, unlike the usual Sweet-Parker result, the width of this inner zone does not determine the reconnection speed. Instead its width automatically adjusts itself to match the reconnection speed determined by ambipolar diffusion. Ambipolar diffusion removes matter from the reconnection zone and enhances the reconnection velocity.

The existence of this broad outflow raises the possibility of a problem in our model. Since the gas outflow is not confined to the layer where the magnetic field is actually recombining, we cannot assume that the ejected gas will be removed from the downstream flow by the relaxation of the reconnected field lines. Instead, the neutral component of the gas has to diffuse into the wedge of reconnected field lines by overcoming the ion-neutral drag. Since the downstream pressure deficit will help push the neutrals in this direction, this diffusion will occur at a rate $\sim \eta_{ambi}/L_y^2$. In the mean time the transverse motion of the neutrals will be approximately along the field lines, since the ions and neutrals are strongly
coupled through collisions and move together at a speed $\sim v_A$. Consequently, the gas will move a distance $v_A L_y^2/\eta_{ambi} \sim L_x$ while diffusing onto the reconnected field lines. It is a basic condition of this model that the reconnection region is embedded in some larger flow, allowing room for the relaxation of field lines and the escape of ejected ions. Since this distance is no larger than the size of the reconnection region it is a short enough that we can assume that the accumulation of ejected gas does not pose a problem for the continued flow of gas from the reconnection region.

We can estimate the width of the resistive zone by considering the motion of the field lines for $y < L_y$. Within this zone the magnetic field lines speed up as the pressure gradient steepens. Since the bulk velocity remains comparable to $V_{rec}$, while the magnetic field strength plummets, the magnetic flux is carried by the second term on the LHS of equation (11). Consequently, we have

$$V_{rec} B_\infty \approx \eta_{ambi} \partial_y B,$$

for $y \ll L_y$ and where $B_\infty$ is the magnetic field strength far outside the reconnection layer. This has the solution

$$B = B_\infty \left( \frac{3V_{rec} y + C}{\eta_{ambi,\infty}} \right)^{1/3},$$

where $\eta_{ambi,\infty}$ is the ambipolar diffusion coefficient far from the reconnection zone and $C$ is a constant equal to $B(y = 0_+)^3/B_\infty^3$. (The magnetic field strength near $y = 0$, but outside the layer where reconnection actually occurs, is defined as $B(0_+)$.)

The ion velocity, which is also the inward speed of the magnetic field lines, is given by

$$v_i = \eta_{ambi} \partial_y \ln B = V_{rec} \left( \frac{3V_{rec} y + C}{\eta_{ambi,\infty}} \right)^{-1/3}.$$

If $C = 0$ then this expression diverges near $y = 0$. One resolution of this problem is to note that $v_i$ is limited by the local value of $v_A/x^{1/2}$, where $x$ is the ionization fraction of the gas. As $v_i$ approaches this value the assumption that the acceleration term is negligible is violated and we have $v_i \rightarrow v_A/x^{1/2}$. This implies a limit on $C$, which can be used as an estimate for $B(y = 0_+)$. Using equations (15) and (16) we find

$$B(0_+) \sim B_\infty \left( \frac{x \eta_{ambi,\infty}}{L_x v_{A,\infty}} \right)^{1/4} = B_\infty \left( \frac{x v_{A,\infty} L_{i,n}}{L_x} \right)^{1/4}.$$

This expression will be relevant provided that $\eta$ remains less than $\eta_{ambi}$ as the magnetic field approaches this asymptotic limit. More precisely, this is relevant when

$$\eta < \eta_{ambi,\infty} \left( \frac{x \eta_{ambi,\infty}}{L_x v_{A,\infty}} \right)^{1/2}.$$
In this case the asymptotic value of \( v_i \) is
\[
v_i(0+) = v_{A,\infty} \left( \frac{v_{A,\infty} t_{n,i}}{x L_x} \right)^{1/4},
\]
and the width of the resistive reconnection region is
\[
\Delta \approx \frac{\eta}{v_i(0+)} \approx \left( \frac{L_x \eta}{v_{A,\infty}} \right)^{1/2} \left( \frac{\eta}{\eta_{ambi,\infty}} \right)^{1/4} \left( \frac{\eta}{L_x v_{A,\infty}} \right)^{1/4}.
\]

In either case, we expect that \( B(0+) \ll B_\infty \) and consequently that magnetic pressure scales as \( y^{2/3} \) throughout the outer layers of the reconnection region. This implies that the overpressure in the layer \( L_y \) scales as \( B_\infty^2 (1 - (y/L_y)^{2/3}) \) and that the outflow velocity will be of the order of \( v_A \). This justifies \textit{a posteriori} our assumption of rapid outflow in this case.

When the condition expressed in equation (18) is not satisfied \( v_i \) stays below the ion Alfvén speed at all times. In this case we can take \( C = 0 \) and determine the width of the resistive reconnection region from the condition that
\[
\frac{\eta}{\Delta} \approx |v_i(y = \pm \Delta)|.
\]
From equation (16) this implies that
\[
\Delta \approx \frac{L_x \eta}{v_{A,\infty}} \frac{\eta}{\eta_{ambi,\infty}}.
\]

We see from equations (20) and (22) that the width of the resistive reconnection region is much smaller in this case than when ambipolar diffusion is negligible. This is expected, since a narrow reconnection region is necessary for a fast reconnection speed, but it also raises the question of local heating. The bulk of the magnetic energy is dissipated through expansion as the ions speed up, there is an irreducible minimum which is annihilated inside the reconnection region proper. The volume heating rate will be
\[
\dot{E} \approx B(y = 0)^2 \frac{\eta}{\Delta^2},
\]
but since \( \Delta = \eta/v_i \) this is
\[
\dot{E} \approx B(y = 0)^2 \frac{v_i^2}{\eta} = B_\infty^2 \frac{V_{rec}^2}{\eta} = B_\infty^2 \frac{v_{A,\infty} \eta_{ambi,\infty}}{L_x \eta}.
\]
a model which relies on an abundance of neutral particles to carry mass away from the reconnection region. We will not attempt to calculate the consequences of this heating here. It is probably not the most severe constraint on ambipolar diffusion in the interstellar medium. Instead, we suggest that the long recombination time for ions, coupled to the effects of ion pressure, poses a much greater problem. We examine this point in the next section of this paper.

3. The Role of Ion Pressure

The critical assumption in the preceding section is that it is reasonable to assume that the ions and neutral particles maintain their usual ratio within the reconnection layer. We can see why this important by restricting our attention to the resistive layer. The net inward flux of ions will be \( > V_{\text{rec}} n_{i,\infty} \). This constitutes a lower limit, since within the whole reconnection region of width \( L_y \) the magnetic field lines speed up, so that their density drops, and the same effect will depress \( n_i \) below its equilibrium value. A uniform neutral distribution will therefore add incoming ions to this flow.

The rate at which ions leave the resistive region through recombination is

\[
n_i(0) \frac{\Delta}{\tau_{\text{recomb}}(0)} = n_i(0) \frac{\eta}{v_i(0_+)^{\tau_{\text{recomb}}(0)}} > n_{i,\infty} V_{\text{rec}} .
\]

(25)

This assumes that the loss of ions through their expulsion in the \( \hat{x} \) direction is negligible, but if that process dominates then we will recover the usual Sweet-Parker result. When the recombination rate is small, ions will accumulate in the resistive layer creating a strong ion pressure gradient, and the reconnection process will be limited by the rate at which ions can recombine with electrons within the resistive layer.

We can understand the role of the ion pressure by equating the ion density flux with the recombination losses in the resistive layer. We have

\[
v_i(0_+) \rho_i(0_+) = \frac{2 \rho_i(0) \Delta}{\tau_{\text{recomb}}(0)} .
\]

(26)

This implies that

\[
v_i(0_+) = \left( \frac{2 \eta(0)}{\tau_{\text{recomb}}(0)} \right)^{1/2} \left( \frac{T(0_+)}{T(0)} + \frac{V_{\text{rec}}}{v_i(0_+)} \frac{B_\infty^2}{8 \pi \rho_{i,\infty} c_s^2(0)} \right)^{1/2},
\]

(27)

where we have invoked pressure balance for the ion density in the resistive layer, and assumed that \( B/\rho_i \) is a constant outside of the resistive layer.
The recombination rate is $\alpha n_e$, where $n_e \approx n_i$ and $\alpha$ is the recombination coefficient (see Spitzer 1978, p. 107), which scales approximately as $2 \times 10^{-11} \phi(T)/T^{1/2}$ cm$^3$ s$^{-1}$ for a hydrogen plasma, where $\phi$ is a slowly varying function of temperature. In an ionized plasma the resistivity is $\approx 2.65 \times 10^{12} T^{-3/2}$, and this expression should remain valid as long as the scattering length for electrons is determined by inelastic collisions with ions. In a largely neutral gas the exact dependence on density and temperature can be fairly complicated. However, for our purposes a crude estimate is all that is required and in any case the density of ions in the resistive layer is greatly enhanced in order to balance the magnetic pressure in the surrounding plasma. Using these results we can rewrite equation (27) as

$$v_i(0_+) = \left( \frac{2\eta(0_+)}{\tau_{\text{recomb}}(0_+)} \right)^{1/2} \frac{T(0_+)}{T(0)}^{3/2} \left( 1 + \frac{V_{\text{rec}}}{v_i(0_+)} \frac{B_\infty^2}{8\pi \rho_{i,\infty} c_i^2(0_+)} \right)^{1/2} \frac{T(0_+)}{T(0)}^{3/2} \left( 1 + \frac{V_{\text{rec}}}{v_i(0_+)} \frac{B_\infty^2}{8\pi \rho_{i,\infty} c_i^2(0_+)} \right) \right),$$

Equation (28) is only an upper limit on $V_{\text{rec}}$, which in principle could be much lower if there exists an outer layer dominated by ion-neutral drag. However, we can show that this is unlikely to be the case. If such a layer exists then we can see from equation (16) that the ion velocity in this layer becomes

$$v_i = V_{\text{rec}} \left( \frac{3V_{\text{rec}}}{\eta_{\text{ambi,}} v_i(0_+)} \right) \left( 1 + \left( \frac{V_{\text{rec}}}{v_i(0_+)} \right)^{3} \right)^{-1/3}.$$

This in turn implies an outer layer width of

$$L_y \sim \frac{\eta_{\text{ambi,}} V_{\text{rec}}^2}{v_i(0_+)^3},$$

and a reconnection velocity of

$$V_{\text{rec}} = \left( \frac{v_i(0_+)^2 L_x}{v_A \eta_{\text{ambi,}} v_i(0_+)} \right) v_i(0_+).$$

This could represent an enormous reduction of $V_{\text{rec}}$ from its upper limit of $v_i(0_+)$, but this outer layer will not form unless the cumulative drag from the neutrals in this layer is large enough to lead to a significant reduction in the local magnetic field strength. Otherwise the outer layer will collapse and $v_i(0_+) \rightarrow V_{\text{rec}}$.

At large distances the ions and neutrals move together with a velocity $V_{\text{rec}}$ towards the reconnection zone. The accumulation of neutrals will lead them to decelerate and accumulate in a layer of width $L_n$. If we assume that this layer is much thicker than the actual zone of reconnection, then the relative ion-neutral velocity in this layer will be $\sim V_{\text{rec}}$ and the net drag force will be

$$\rho_n V_{\text{rec}} \frac{L_n}{t_{n,i}} \sim \delta P_n,$$

where $\delta P_n$ is the pressure difference.
where we have assumed that the drag force is balanced by the excess neutral pressure. The neutrals will be expelled laterally at a speed $v_{x,n}$ given by

$$
v_{x,n}^2 \sim \frac{\delta P_n}{\rho_n}.
$$

(33)

Using Eqs. (32), (33) and conservation of mass we conclude that

$$
v_{x,n} \sim \left(\frac{v_{rec}^2 L_x}{t_{n,i}}\right)^{1/3},
$$

(34)

and

$$
L_n \sim L_x \left(\frac{v_{rec} t_{n,i}}{L_x}\right)^{1/3}.\quad (35)
$$

The condition that the total drag within this layer has a negligible effect on the reconnection layer is $\delta P_n/\rho_n \ll v_A^2$ or $v_{x,n} \ll v_A$ (which is also the condition that the neutral ejection velocity be much less than $v_A$). This is equivalent to

$$
V_{rec}^2 \ll \frac{\eta_{ambi} v_A}{L_x}.
$$

(36)

In other words, as long as the reconnection speed is less than we would obtain from the naive substitution of $\eta_{ambi}$ for $\eta$ in the usual Sweet-Parker formula, we can ignore the ion-neutral drag outside the reconnection layer and set $v_i(0_+) = V_{rec}$ in equation (28).

A significant complication is that the momentum of the ejected neutrals may diffuse over a distance greater than the value of $L_n$ given in equation (35), thereby increasing $L_n$ and possibly the role of the neutral drag in the reconnection layer. This effect will become important when

$$
\frac{\nu_n}{L_n^2} \sim \frac{v_{x,n}}{L_x},
$$

(37)

where $\nu_n \sim c_s^2 t_{n,n}$ is the neutral gas viscosity. In this limit we need to replace equation (33) with

$$
\frac{\delta P_n}{L_x} \sim \rho_n \frac{\nu_n}{L_n^2} v_{x,n}.
$$

(38)

If we combine this result with equation (32) and the condition that $V_{rec} L_x \sim v_{x,n} L_n$ we get

$$
L_n \sim \left(c_s^2 t_{n,n} t_{n,i}\right)^{1/4} L_x^{1/2}.
$$

(39)

The condition that the ion-neutral drag does not dissipate a significant fraction of the magnetic energy outside the reconnection layer is $v_A^2 \gg V_{rec} L_n/t_{n,i}$, which can be rewritten, with the aid of equation (39) as

$$
v_A^2 \gg V_{rec} \left(\frac{L_x^2 v_{x,n}^2}{t_{n,i}}\right)^{1/4},
$$

(40)
or, using the definitions of $V_{\text{max}}$ and $\eta_{\text{ambi}}$,

$$V_{\text{rec}} \ll V_{\text{max}} \left( \frac{v_A}{c_s} \right)^{1/2} \left( \frac{t_{n,i}}{t_{n,n}} \right)^{1/4}. \tag{41}$$

For largely neutral gases the ratio of $t_{n,i}$ to $t_{n,n}$ will be of order $x^{-1}$, and $v_A$ is usually of order $c_s$. Consequently, under typical conditions in the interstellar medium equation (41) will be satisfied by a comfortable margin (cf. Table 1).

Once again we need to consider whether or not the broad, slow outflow of neutrals from the reconnection region will result in an accumulation of gas outside the reconnection region, along magnetic field lines that have not yet undergone reconnection. The appropriate diffusion coefficient is, as before, $v_A^2 t_{n,i}$. The diffusion scale is $L_n$, the thickness of the stagnation layer in the neutral flow. We see from equations (42), (34), and (35) that this implies that the transverse distance covered by the neutrals while diffusing into the reconnected wedge will be

$$v_{x,n} L_n^2 \eta_{\text{ambi}} \approx L_x \left( \frac{V_{\text{rec}}}{V_{\text{max}}} \right)^{4/3}. \tag{42}$$

Since this is always less than $L_x$ we conclude that the ejected neutrals can be ignored.

This case is illustrated schematically in figure 1, with the local velocities of the neutrals and ions indicated by thin and thick arrows. The material in the reconnection layer is actually moving in the $\hat{x}$ direction with a speed $\sim v_A$, but since most of the gas leaves the reconnection layer in the $\hat{y}$ direction as neutral particles we have ignored the motion of the ions in that layer.

Our only remaining worry is that the transverse expulsion of neutrals may serve to remove large numbers of ions from the reconnection zone. It is certainly reasonable to assume a tight coupling between the ion and neutral transverse speeds, since the transverse shear rate, $v_{x,n}/L_x$ is much smaller than the ion coupling rate, $(\rho_n/\rho_i)t_{n,i}^{-1}$. However, for typical ISM parameters we also have $\tau_{\text{recomb}}^{-1} > v_A/L_x > v_{x,n}/L_x$, so the ion fraction in the gas is maintained even while individual ions are expelled. We also note that even if we ignored recombination the equations of continuity for ions and neutrals combine to give

$$v_i \partial_y \ln \rho_i - v_n \partial_y \ln \rho_n + \nabla (v_i - v_n) = 0. \tag{43}$$

Within the outer layer of thickness $L_n$ the neutral particles come to a halt and $v_i$ remains close to $V_{\text{rec}}$. Consequently the density scale height for the ions is not less than $L_n$ and some large fraction of the ions will reach the resistive layer, despite transverse losses.

We conclude that

$$V_{\text{rec}} = \left( \frac{2 \gamma_{\infty}}{\tau_{\text{recomb,}\infty}} \right)^{1/2} \left( \frac{T_{\infty}}{T(0)} \right)^{3/2} \left( 1 + \frac{1}{\beta x} \right) \tag{44},$$
where $\beta$ is the ratio of the gas pressure to the magnetic pressure. The width of the reconnection layer is

$$\Delta = \frac{\eta(0)}{V_{\text{rec}}} = \left(\frac{\eta_\infty \tau_{\text{recomb},\infty}}{1 + \frac{1}{\beta x}}\right)^{1/2}.$$  

(45)

The principle source of ion heating in the reconnection layer is ohmic heating due to the dissipation of the magnetic field. Cooling can take place either through radiative losses or through the transfer of energy to the neutral gas. The neutrals will, in turn, shed their excess thermal energy through radiative losses and/or by escaping from the reconnection layer and diluting the extra heat over a large volume. Since neutral mean free path is typically much larger than $\Delta$, and since $c_s \gg \Delta/t_{i,n}$, the ion cooling rate will be controlled by the ion-neutral collision rate. In other words,

$$\frac{B_\infty^2 \eta(0)}{8\pi \Delta^2} \approx \frac{3}{2} 2 n_i k_B T(0) \frac{2}{t_{i,n}},$$

(46)

where we have assumed that $T(0) \gg T_\infty$, since otherwise the correction is of no interest, and used $B_\infty$ since we have already shown that $B_\infty \sim B(0_+)$ for reconnection in the interstellar medium. We have also assumed that the electrons and ion share the same temperature in the reconnection layer. Given that the charged particle pressure in the ionization zone balances the external magnetic pressure, equation (46) implies

$$\frac{V_{\text{rec}}^2}{\eta(0)} \approx \frac{3}{t_{i,n}}.$$  

(47)

We can rewrite this with the aid of equation (44) as

$$\left(1 + \frac{1}{\beta x}\right)^2 \approx \left(\frac{3 \tau_{\text{recomb},\infty}}{t_{i,n}}\right) \frac{\phi(T_\infty)}{\phi(T(0))} \left(\frac{T(0)}{T_\infty}\right)^{5/2}.$$  

(48)

This implies that when reconnection layer heating is important we should substitute the expression

$$\left(\frac{T_\infty}{T(0)}\right)^{3/2} \approx 31 \left(1 + \frac{1}{\beta x}\right)^{-6/5} x^{-3/5} \phi(T(0))^{-3/5} T_\infty^{0.3},$$  

(49)

into equation (46). We see at once that $T(0) > T_\infty$, so that this correction is appropriate, only in cold regions with $x \ll 10^{-3}$ (assuming $\beta$ is of order unity). As we reach this limit we go from $V_{\text{rec}} \propto x^{-0.5}$ to $V_{\text{rec}} \propto x^{0.1}$.

It is important to remember that for strictly one dimensional formulations of this problem there is no natural choice for $L_n$ and the actual reconnection speed will remain
sensitive to the computational box size and/or the boundary conditions. Although \( L_x \) does not appear in equation (44) it is still present implicitly as part of the constraints that make this a well-posed problem. Purely one dimensional calculations will not capture all the relevant physics, and may result in a wide variety of estimates for \( V_{\text{rec}} \) and \( \Delta \).

Equation (44) constitutes the main result of this paper, and may be regarded as a generalization of the Sweet-Parker reconnection rate to an ambipolar medium with a long recombination time. The major effect we have neglected here is additional ionization within the reconnection layer caused by the reconnection process. We have included the effect of local heating on the resistivity, gas density, and recombination rate. We note that even taking \( T(0) = T_\infty \) we can show that reconnection in this model is rather slow.

We should note that equation (44) is correct only if \( \tau_{\text{recomb}} \) is much less than the shearing time \( L_x/v_A \). In the opposite limit, ions will escape from the sides of reconnection zone rather than through recombination. In this case we are almost back in the regime described in section 2.1. However, there are two important physical effects which have to be considered. First, we have already seen that the partial decoupling of ions and neutrals implies a large concentration of ions in the reconnection layer itself. Consequently the rate of ion ejection is enhanced by the ratio \( n_i(0)/n_{i,\infty} \). The ion conservation equation becomes

\[
\Delta v_{\text{eject}} n_i(0) = L_x V_{\text{rec}} n_{i,\infty}.
\]

Second, the neutrals will diffuse out the reconnection layer, spreading the transverse momentum and creating an additional drag on the ions. Since these effects work in opposite directions, it’s not immediately obvious whether the presence of neutrals increases or decreases the reconnection speed when \( \tau_{\text{recomb}} \) is very large.

Balancing the ion pressure gradient in the \( \hat{x} \) direction with the neutral drag, we get

\[
\frac{\delta P_i}{L_x} \approx \rho_i(0) \frac{\Delta v_x}{t_{i,n}}.
\]

The transfer of momentum to the neutrals is then balanced by the viscous drag on the neutrals, so that

\[
\rho_i(0) \frac{\Delta v_x}{t_{i,n}} \Delta \approx \rho_n v_{\text{eject}} \frac{v_n}{L_n},
\]

where we have assumed that \( \Delta v_x \), the difference between the ion and neutral particle velocities in the \( \hat{x} \) direction, is small compared to \( v_{\text{eject}} \). The width of the spread of \( \hat{x} \) momentum is given by equation (37), so we conclude that

\[
v_A^2 \Delta \approx \left( v_{\text{eject}}^3 L_x v_n \right)^{1/2}.
\]
Remembering that $\Delta \approx \eta(0)/V_{\text{rec}}$ we can combine equations (50) and (53) to obtain

$$V_{\text{rec}} \approx \left( \frac{v_A \eta(0)}{L_x} \right)^{1/2} \left( \frac{\eta(0)}{\nu_n} \right)^{1/8} \left( \frac{n_i(0)}{n_i,\infty} \right)^{3/8}.$$

(54)

Consequently, we can estimate the corrected Sweet-Parker reconnection velocity as

$$V_{\text{CSP}} \approx \left( \frac{v_A \eta_\infty}{L_x} \right)^{1/2} \left( \frac{\eta_\infty}{\nu_n} \right)^{1/8} \left( \frac{T_\infty}{T(0)} \right)^{21/16} \left( 1 + \frac{1}{x\beta} \right)^{3/8}.$$

(55)

In spite of the fact that typically $\nu_n \gg \eta_\infty$, this estimate will usually be larger than the standard SP reconnection speed for largely neutral, magnetized gas.

Finally, we note that our estimates are based on the assumption that the reconnection layer is in a steady state, meaning that the local pressure excess is comparable to the magnetic field pressure and that the accumulation of ions in the reconnection layer has reached a steady state. If the magnetic field lines are being compressed by turbulent forces then the former condition will be reached in about one eddy turnover time, or, assuming $v_A \sim v_{\text{turb}}$, about $L_x/v_A$. Otherwise we should replace the eddy turnover time with the large scale time dynamical time scale. At earlier times the confining pressure will be less than our estimates, and the reconnection rate slower. On the other hand, filling up the reconnection layer with ions will take a time $\sim (\Delta/V_{\text{rec}})(n_i(0)/n_i,\infty)$ and at earlier times the reconnection layer will be compressed to a smaller width, with a consequent increase in the reconnection rate. We can rewrite this time scale using equation (44). We obtain

$$\frac{\Delta}{V_{\text{rec}}} \frac{n_i(0)}{n_i,\infty} \approx \frac{\eta(0)}{V_{\text{rec}}^2} \frac{n_i(0)}{n_i,\infty} \approx \tau_{\text{recomb},\infty} \left( \frac{T(0)}{T_\infty} \right)^{1/2} \left( 1 + \frac{1}{x\beta} \right)^{-1}.$$

(56)

This will usually be somewhat less than an eddy turnover time in the interstellar medium.

We conclude that the early phases of reconnection are apt to be characterized by weak compression and slow reconnection, and in any case that the reconnection of magnetic flux is unlikely to proceed much more rapidly in the early phases of reconnection.

4. Phases of the ISM and reconnection

Does ambipolar diffusion actually lead to a dramatic enhancement of reconnection rates in the ISM? In Table 1 we give our results for idealized partially ionized phases of the ISM. We include, for reference, $V_{\text{max}}$, the reconnection velocity when $\tau_{\text{recomb}} \to 0$, and $V_{\text{SP}}$, the Sweet-Parker reconnection speed without corrections for natural drag or ion density enhancement. More realistically, we also show $V_{\text{rec}}$, in which the reconnection speed
is limited by the recombination rate, and \( V_{\text{CSP}} \), the corrected Sweet-Parker reconnection speed. For molecular hydrogen clouds the reconnection layer heating is important in \( V_{\text{rec}} \), and has been included using equation (49). The reconnection velocities \( V_{\text{SP}} \) and \( V_{\text{CSP}} \) will depend on the assumed values of \( L_x \). Here we take \( L_x = 1 \) pc so that both these speeds scale as \((1 \text{ pc}/L_x)^{1/2}\). We also make use of the evidence that for all the phases shown in Table 1 the Alfvén velocity \( v_A \) is of the order of a couple of kilometers per second. Here we assume that it is \( 10^5 \) cm s\(^{-1}\). Finally, each ISM phase we give \( n_f \), the mass fraction of that phase. A discussion of the values we use here is contained in Draine and Lazarian (1997). We note that an additional 22% is contained in ionized gas.

There are several important points to be noted in connection with Table 1. First for all the phases presented in Table 1 it is the “recombination limited” reconnection rate (see Eq (44)) that determines the reconnection speed on scales of \( \sim 1 \) pc. Apart from diffuse warm HI, where reconnection is slow anyhow, the corrected Sweet-Parker reconnection speed given by equation (55) is at least one order of magnitude slower than \( V_{\text{rec}} \). This situation changes only at much smaller scales where outflows from reconnection regions can carry out most of the ions faster than recombination can eliminate them.

Second, in all phases of the ISM, \( V_{\text{max}} \) is clearly a gross exaggeration of the degree to which ambipolar diffusion can enhance reconnection. Nevertheless, \( V_{\text{rec}} \) is very much larger than \( V_{\text{CSP}} \) for all the cold phases of the ISM. The latter is, in turn, considerably larger than \( V_{\text{SP}} \). In both cases, this increase is due to the compression of ions in the reconnection layer. The conclusion that \( V_{\text{rec}} > V_{\text{CSP}} \) is due to the double contribution of the ion density enhancement in the former case, that is, not only does it contribute directly by enhancing the efficiency of ion loss relative to the incoming flux of ions, but it also contributes by raising the recombination rate. This explains why the warm neutral gas shows only a small increase in its reconnection speed. It also suggests that weak magnetic fields will have much slower reconnection speeds.

Third, \( V_{\text{rec}} \) is always much smaller than \( v_A \). The only case in which it exceeds \( 10^{-3} \times v_A \) is that of dense molecular gas. Moreover, this phase has a modest mass fraction of about 3%. While this result implies that reconnection should be efficient enough to remove magnetic flux from star forming regions, and disconnect the internal magnetic field from its environment, it does not explain how reconnection could be fast enough to sustain the galactic dynamo.

Finally, these calculations are based on the assumption that the plasma can be treated as a fluid. Given the low collision rates in the ISM, this assumption requires, at a minimum, that all scales of interest be much greater than the ion Larmor radius. In particular, we
require that $\Delta \gg R_{L,\text{ion}}$ or

$$1 \gg \frac{V_{\text{rec}} V_{\text{th,ion}} m_{\text{ion}} c}{\eta(0) e B} \approx \left(\frac{2}{\eta(0) \tau_{\text{recomb},\infty}}\right)^{1/2} \left(1 + \frac{1}{\beta x}\right) \frac{V_{\text{th,ion}} c x^{1/2}}{v_A \omega_{p,\text{ion}}},$$

(57)

where $\omega_{p,\text{ion}}$ is the ion plasma frequency. If we evaluate this expression for the idealized ISM phases shown in Table 1 we find that it rises from a low of about $3 \times 10^{-2}$ for diffuse cold HI to about 0.13 for diffuse molecular gas. There is a weak dependence on reconnection layer heating which increases the ratio for molecular gas. As before, we have calculated the local heating using equation (49). For dense molecular gas the ion Larmor radius is roughly the same as the reconnection layer thickness. This suggests that it would be appropriate to reconsider this case without the use of the fluid approximation. Such a calculation is beyond the scope of this paper. We merely note that our results for this case should be treated with caution. However, it seems unlikely that the reconnection rates will decrease significantly when we allow for the finite Larmor radius of the ions. If $r_{L,\text{ion}} \gg \Delta$ then the ions do not feel the magnetic fields in the reconnection layer. Nevertheless, they will be confined to this layer by the electric field of the electrons, which in turn confined by the magnetic pressure exerted on the current-carrying component of the plasma. It follows that the electrons and ions will be moving together as in the hydrodynamic approximation. Possible deviations from this approximation in the form of instabilities, if they take place, will only enhance magnetic diffusivity and reconnection.

5. Discussion and Conclusions

In this paper we have examined the role of neutral particles on the structure of reconnection regions in the ISM. In so doing, we have assumed that the basic geometry posited by Sweet and Parker accurately describes reconnection in three dimensions. We find that ambipolar diffusion can give rise to much faster reconnection speeds than those obtained by the naive use of the usual Sweet-Parker formula, but these speeds are still very small compared to the local Alfvén speed. In addition, most of the increase is due to the compression of ions in the reconnection layer, and would not occur in a fluid where the ion and magnetic pressures were comparable.

How can we interpret these results? Obviously, there are physical effects which we have ignored here, but which might affect the geometry or the physics of reconnection. Assuming our results are a realistic description of reconnection in the ISM we can examine their implications for the structure of the magnetic field in the ISM. We will address the latter issue first, and then briefly take up the prospect of obtaining faster reconnection speeds under realistic circumstances.
What have we learned about flux tube formation? In a previous paper (Vishniac 1995) one of us showed that flux tubes should be a natural feature of turbulent high beta plasmas. This is consistent with observations of the solar photosphere, but not otherwise directly testable. In parts of the interstellar medium the effective plasma beta can be large, if one accounts not only for thermal plasma pressure, but for the pressure from compressional shocks and cosmic rays (Lazarian & Vishniac 1996). Therefore, it seems possible that flux tubes could form in the interstellar medium. The basic mechanism involves the process of turbulent pumping, in which flux tubes are stretched by the surrounding turbulence, and then twisted to produce close loops of flux which self-destruct and release their entrained ions into the surrounding medium. This assumes that reconnection is efficient.

However, our results indicate that ambipolar diffusion acts against the flux tube formation process. In addition to the small reconnection speeds we have obtained, ambipolar diffusion makes the tubes leaky since it allows neutrals to cross magnetic field lines, diffuse into the flux tubes, and subsequently fill them up with newly produced ions. When recombination times are short we can show that flux tubes can still be formed. In this case ambipolar diffusivity works in exactly the same way that Ohmic resistivity works in a completely ionized plasma. The trade-off between more rapid ambipolar reconnection and ambipolar leakage results in a rather weak dependence of turbulent pumping efficiency on the ambipolar diffusivity. The situation changes dramatically if the recombination times are large, as they are in the ISM. In this case ambipolar reconnection alone cannot provide the necessary rates of turbulent pumping.

We conclude that flux tubes will not form in the ISM unless there are ways to obtain dramatically higher reconnection speeds. The only possible exceptions are the dense molecular parts of the ISM. On the other hand, we note that the best evidence for the existence of flux tubes comes from the solar photosphere, where the high collision rates imply that ambipolar diffusion is only slightly more effective than ohmic diffusion. In addition, the best evidence for high reconnection speeds comes from parts of the solar chromosphere and corona where the neutral fraction is negligible. Evidently there must be circumstances where reconnection is much faster than our estimates and these may allow the formation of flux tubes in the ISM.

What do our results say about the galactic dynamo? Our estimates for $V_{CSP}$ depend on $L_x$, and reconnection might occur over small distances at speeds higher than those we have given here. On the other hand, we have already assumed $L_x$ is of order a parsec, and the large scale eddies in the ISM may be as much as two orders of magnitude larger. Moreover, it is precisely these large scale motions which are usually thought to drive the galactic dynamo. Finally, we note that $V_{rec}$ is insensitive to $L_x$. While it is difficult to
be precise about the reconnection speeds necessary to allow turbulent diffusion and the existence of a large scale dynamo, the usual assumption is that these speeds must be in the range 0.01 to 0.1 times $v_A$ (see Ruzmaikin, Sokoloff & Shukurov 1988) or the magnetic field will become tangled on the smallest scales and show little evidence for large scale organization. Our largest estimates for $v_A$ falls just outside this range and applies to only a few percent of the interstellar gas. Once again our results are disappointing in this regard and suggest the necessity for some way to obtain much larger reconnection speeds.

We could consider circumstances where the gas density is much larger, for example, a cool stellar atmosphere. The gas in the solar photosphere is largely neutral and dense enough that recombination will be very fast. On the other hand, under these conditions $t_{i,n}$ is also very much shorter, and ambipolar diffusion is not very much faster than ohmic diffusion. For example, if we take $n_n \sim 10^{10}$, and $B \sim 1$ gauss then $\eta_{ambi} \sim 10^{11}$, about $10^2$ larger than the ohmic diffusivity in a cool stellar atmosphere. This suggests a reconnection speed an order of magnitude larger than the Sweet-Parker estimate, but still quite small compared to turbulent diffusivity in a giant stellar atmosphere.

Are there additional physical effects that might affect our results? We start by considering cosmic rays. Cosmic rays are an essential component of the interstellar medium and their pressure is roughly of the same order as the pressure of the magnetic field in the ISM, suggesting that they may play an important role in all plasma processes in the ISM, including reconnection. Our discussion here will follow Longair (1994).

Cosmic rays diffuse through magnetic field and the magnetic fluctuations on the scale of the Larmor radius are the most effective source of scattering. Therefore for our order of magnitude estimates we may consider that the time for scattering of cosmic rays is of the order $t \sim N/\omega_L$, where $\omega_L$ is the Larmor frequency and $N$ can be defined as the number of orbits required to substantially alter the initial motion of the charged particle. In other words, if in the course of an individual scattering event the particle deviates over an angle $\Phi \sim B_{rms}/B_{reg}$, where $B_{rms}$ and $B_{reg}$ are the rms and regular components of the magnetic field at the place of scattering, then $N$ satisfies the relation: $N^{1/2} \Phi \sim 1$, where the exponent 1/2 appears above due to the random walk nature of the process. As a result the diffusion coefficient for a charged particle with elementary charge $e$ and energy $E$ can be defined as:

$$\eta_{cr} \sim \frac{r_L^2}{t} \sim \frac{cE}{eB} \Phi^2 . \quad (58)$$

In the zone of reconnection cosmic rays are streaming along the magnetic field lines and perturb magnetic field lines through nonlinear interactions. In these circumstances $\Phi$ does not differ much from unity. As a result values of $\eta_{cr}$ can be as large as $10^{20}$ cm$^2$ s$^{-1}$ for the Mev charged particles. If we assume that the background gas density is negligible, as it
may be in the galactic halo, then the Alfvén speed approaches that of light, and we obtain high reconnection rates (Parker 1992).

In the galactic disk, where the density of gas is not negligible, the situation is quite different. In general the Larmor radius for cosmic ray particles will be much larger than the reconnection layer width. Consequently, cosmic rays cross the reconnection zone in a single orbit. Moreover, neutral particles damp the Alfvénic waves, thereby reducing the coupling between cosmic ray streaming and perturbations to the local magnetic field. Finally, the reduced magnetic field in the reconnection layer decreases the coupling between the plasma and the cosmic rays. In the end, the only obvious role for the cosmic rays is that they increase the effective pressure outside the reconnection zone, and consequently help compress the gas in the reconnection layer. This changes the relation between the mass influx and mass efflux so that:

\[
\rho_1 L_x u_x = \rho_2 L_y u_y ,
\]

(59)

where

\[
\frac{\rho_2}{\rho_1} = \frac{P_{\text{gas}}}{P_{\text{gas}} + P_{\text{cr}}} .
\]

(60)

Above \(P_{\text{gas}}\) and \(P_{\text{cr}}\) are the gaseous and cosmic ray pressure respectively. This increases the reconnection speed by the square root of the density ratio. Since the cosmic ray pressure is roughly the same order as the gas pressure in the interstellar medium, this implies that cosmic rays change the reconnection rate just by a factor of order unity.

How robust is our assumption concerning the basic reconnection geometry? It is well known that the kind of narrow current sheet we are positing as the zone of reconnection is dynamically unstable to tearing modes (Furth, Killeen, & Rosenbluth 1963). This instability persists for reconnection in largely neutral environments, as noted by Brandenburg and Zweibel (1995). Turbulent mixing of the current sheets will promote reconnection by enhancing the transport of fresh magnetic flux into the reconnection zone and by mixing the accumulating ions outward. In fact, at least one author (Strauss 1988) has proposed that this will lead to reconnection speeds of order \((\delta B/B)v_A\) and shown that \(\delta B/B\) will be of order unity in saturated turbulence. Although some enhancement of reconnection speeds seems inevitable, there are grounds for doubting that this will dramatically change our results. We will defer a detailed discussion until our next paper, but preliminary estimates indicate that allowing for the effects of three dimensional nonlinear tearing modes in an ionized medium will increase the reconnection speed over the Sweet-Parker estimate by only a factor of the magnetic Reynolds number to the one sixth power.

Finally, and on a more hopeful note, the existence of a third dimension may allow the interpenetration of magnetic field lines for favorable geometries, i.e. an interchange
instability, which will facilitate reconnection by dramatically increasing the surface area of the reconnection layer. (cf. Uchida & Sakurai 1977). We will explore these ideas in a subsequent paper.

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Table 1. Reconnection speeds for idealized phases of the ISM

| ISM phase               | $\tau_{\text{recomb}}$ | $V_{\text{max}}$ | $V_{\text{rec}}$ | $V_{SP}$  | $V_{CSP}$ |
|-------------------------|-------------------------|-------------------|-------------------|-----------|-----------|
| diffuse cold HI         | $4.0 \times 10^{12}$    | $2.3 \times 10^3$ | 7.8               | $9.3 \times 10^{-3}$ | $1.4 \times 10^{-1}$ |
| $n = 30 \text{ cm}^{-3}$, $x = 10^{-3}$ | $T = 100 \text{ K}$, $n f \approx 0.15$ |                   |                   |           |           |
| diffuse warm HI         | $3.6 \times 10^{13}$    | $1.9 \times 10^3$ | $4.1 \times 10^{-4}$ | $4.3 \times 10^{-4}$ | $3.0 \times 10^{-4}$ |
| $n = 0.3 \text{ cm}^{-3}$, $x = 0.15$ | $T = 6000 \text{ K}$, $n f \approx 0.4$ |                   |                   |           |           |
| diffuse H$_2$           | $1.6 \times 10^{13}$    | $5.6 \times 10^3$ | (38)              | $1.4 \times 10^{-2}$ | 0.56      |
| $n = 50 \text{ cm}^{-3}$, $x = 10^{-4}$ | $T = 60 \text{ K}$, $n f \approx 0.2$ |                   |                   |           |           |
| dense H$_2$             | $4.5 \times 10^{12}$    | $3.9 \times 10^3$ | (220)             | $3.1 \times 10^{-2}$ | 4.0       |
| $n = 10^3 \text{ cm}^{-3}$, $x = 10^{-5}$ | $T = 20 \text{ K}$, $n f \approx 0.03$ |                   |                   |           |           |

Note. — All quantities are given in cgs units. $\tau_{\text{recomb}}$ is the recombination time outside the reconnection layer; $V_{\text{max}}$ follows from infinitely fast recombination (Eq. (12)); $V_{\text{rec}}$ allows for a realistic recombination rate (Eq. (44)); $V_{SP}$ is the Sweet-Parker reconnection rate (Eq. (1)); and $V_{CSP}$ is the corrected Sweet-Parker reconnection speed given by Eq. (55). $V_{\text{rec}}$ for H$_2$ is placed in brackets as it allows for heating of the reconnection layer.
Fig. 1.— A schematic of a reconnection region. Magnetic reconnection is occurring in a layer of width $\Delta$. Gas velocities are marked with thick lines for ions and thin lines for neutrals. The gas outflow is confined to a region of width $L_n$. The $\hat{x}$ axis is horizontal and the $\hat{y}$ axis is vertical.