SESSION I

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MAGNETIC FIELD RECONNECTION IN COSMIC PLASMAS

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ABSTRACT

A brief review is presented of the concept of magnetic field reconnection or merging. This process occurs whenever an electric field is present along a separator line in the magnetic field. The basic properties of reconnection are discussed in the context of the classical MHD models by Sweet and Parker and by Petschek. Attention is then focussed on reconnection in collision-free plasmas. The energization of charged particles during their interaction with the current layers associated with the reconnection geometry is discussed and the nature of the processes occurring in the so-called diffusion region which surrounds the separator is considered. Finally, comments are made on the nonsteady aspects of reconnection at the earth's magnetopause.

1. INTRODUCTION

Magnetic field reconnection, or merging, is a universal process for the conversion of magnetic energy into plasma kinetic and thermal energy. The process, which may occur either impulsively or in a steady state, taps the magnetic free energy associated with electric current sheets and other sheared field configurations. It is believed to be important in a variety of cosmic situations: solar flares, solar magnetic-field evolution, and perhaps coronal heating; planetary magnetopauses and tails; cometary tails; accretion disks, etc. The process has also been studied extensively in a variety of laboratory devices and simulations as well as in computer simulations.

This paper has two purposes: to provide a brief review of the two classical reconnection models by Sweet-Parker and by Petschek; and to discuss some of the current concerns and findings about reconnection in its magnetospheric setting. For more detailed discussion, the reader is referred to the review papers by Vasyliunas (1975) and Sonnerup (1979).

In the magnetosphere we have a unique opportunity to learn about

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reconnection in a cosmic plasma from direct in situ measurements and it is expected that the information thus obtained can be translated to other cosmic situations, at least to some extent. However, since its introduction into magnetospheric physics by Dungey (1961), the reconnection concept has been a source of continual controversy. Initially, the evidence for reconnection was mostly indirect but the process nevertheless proved to be a powerful organizing concept for a great variety of observations. Recently, more direct evidence for the occurrence of reconnection has become available. But the controversy has not disappeared, for the reconnection process has proved to be far more complicated than originally envisaged. Some of the complexity of the process will become apparent in the following pages.

2. DEFINITION

In order to deal with reconnection in an organized fashion, it is desirable first to provide a simple and unambiguous definition of the process. This definition contains four parts:

(i) A "separatrix" is a magnetic field line surface which separates different magnetic cells as illustrated in Figure 1.

(ii) A "separator" is the line of intersection of a separatrix with itself or perhaps with another separatrix.

(iii) "Reconnection" or "merging" occurs when an electric field $E_{\parallel}$ is present along a separator.

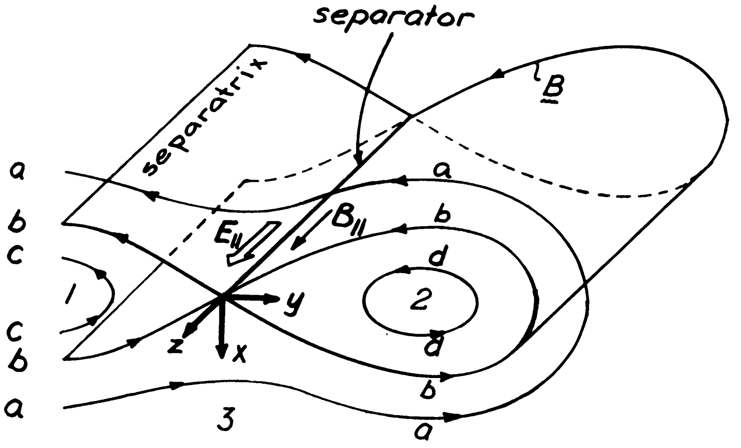

Figure 1. Basic reconnection geometry. The field line "a", originally located in cell 3 moves toward the separatrix surface and lies in that surface at location "b". Reconnection occurs at the separator and the field line is broken into two parts, "c" and "d", located in cells 1 and 2.

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(iv) The "reconnection rate" is given by $E_{\parallel}$ and is equal to the amount of magnetic flux transported per unit time across a unit length of the separator.

Several notes on these definitions should be made:

(a) In general, a separatrix surface is associated with two hyperbolic null points in the magnetic field (see Dungey, 1963). The case shown in Figure 1 is degenerate in the sense that the magnetic field is zero along the entire separator. Usually, this is not the case but a field component $B_{\parallel}$ is present along most of that line.

(b) The term "separator" is synonymous with "reconnection line," "merging line," or "X line." The phrases "neutral line" and "null line" are sometimes used as well but they are somewhat misleading because the field is in general not equal to zero along the separator.

(c) The definition draws only upon the universally accepted concepts of electric and magnetic fields. A number of magnetospheric scientists are uncomfortable with the ideas of moving field lines and frozen magnetic fields so that a definition based on fundamentals only is desirable, even though the term "reconnection" itself invokes the idea of moving field lines. Note that the definition is such that reconnection can occur freely in a vacuum. However, as a process, it is only of interest in the presence of a highly conducting plasma in which the electromagnetic fields and currents are generated self-consistently by the differential motion of ions and electrons.

(d) The definition emphasizes flux transport rather than plasma transport between different topological cells. A well-known mathematical theorem states that $\underline{E}\cdot\hat{\underline{B}} = 0$ is a sufficient condition for the flux transport velocity $\underline{v}_{\phi} = \underline{E}\times\underline{B}/B^2$ to move points that were once joined by a field line in such a manner that they remain joined by a field line at all later times. It is on this fact that the concept of moving field lines is based. Figure 1 illustrates how a field line in cell 3 moves towards the separator. Along this line $\underline{E}\cdot\hat{\underline{B}} \neq 0$ so that the theorem is violated. Reconnection occurs, the result being a transport of magnetic flux from cell 3 to cells 1 and 2. No plasma physics has been introduced into the above discussion but it is the presence of a highly conducting plasma that assures that the condition $\underline{E}\cdot\hat{\underline{B}} = 0$ is satisfied everywhere except at the separator. It is also known that $\underline{E}\times\underline{B}/B^2$ is the electric drift velocity of charged particles. Other drifts such as inertia and gradient drifts are unimportant in most of the external flow but become significant near the separator. It follows from these facts that in a highly conducting plasma, the definition presented here is in all practical respects indistinguishable from the nonlocal definition in terms of "plasma flow across a separatrix" adopted by Vasyliunas (1975).

(e) The amount of magnetic flux transported per unit time across a length element $\underline{d\ell}$ of the separator is

$$d\Phi_m = \underline{B}\cdot\{\underline{v}_\phi \times d\underline{\ell}\} \tag{1}$$

which, upon use of the expression for $\underline{v}_\phi$ simplifies to

$$\frac{d\Phi_m}{d\ell} = E_\parallel \tag{2}$$

It is for this reason that $E_\parallel$ is used as a measure of the reconnection rate. It is however often useful to define a nondimensional rate as

$$M_{A1} \equiv E_\parallel / v_{A1} B_1 \tag{3}$$

where v_{A1} and B_1 are the Alfvén speed and magnetic field at a chosen reference point and reference time. This usage is particularly common in steady-state two-dimensional reconnection models in which $E_\parallel$ can be shown from Faraday's law to be constant throughout the plane perpendicular to the separator. In that case $E_\parallel / B_1$ represents the flux transport velocity (or the electric drift velocity) v_1 at the reference point so that $M_{A1} = v_1/v_{A1}$ is the Alfvén Mach number at that location. The definition remains arbitrary in the sense that the location of the reference point relative to the separator may be chosen differently. This point will be taken up later on.

(f) It is noted that $E_\parallel$ and thus the reconnection rate is invariant under Galileo transformations. In other words, it does not matter whether $E_\parallel$ is measured in a frame of reference in which the separator (or a segment thereof) is at rest or whether it is measured at the instant a moving separator passes the observation point.

(g) Finally, the definition is local in nature and thus cannot and does not distinguish between electrostatic and inductive contributions to $E_\parallel$.

3. SWEET-PARKER MODEL

The Sweet-Parker model of reconnection (Sweet, 1958; Parker, 1963) is shown in Figure 2a. It describes the slow steady-state inflow of two oppositely and strongly magnetized plasmas towards a current sheet and the subsequent rapid outflow of weakly magnetized plasma along the sheet. The basic features of this model can be understood in three simple steps.

First, the pressure at the separator (which is perpendicular to the plane of the figure) is equal to $P_\infty + B_\infty^2/2\mu_o$ where the subscript ∞ denotes upstream conditions. The excess pressure, $B_\infty^2/2\mu_o$, is used to accelerate the outflowing plasma along the sheet so that, in accordance with Bernoulli's law, we have $P_\infty + B_\infty^2/2\mu_o = P_\infty + \frac{1}{2}\rho v_{out}^2$. The pressure at the exit is assumed to be the same as in the inflow and, for simplicity, the density ρ is taken to be a constant. It follows that the

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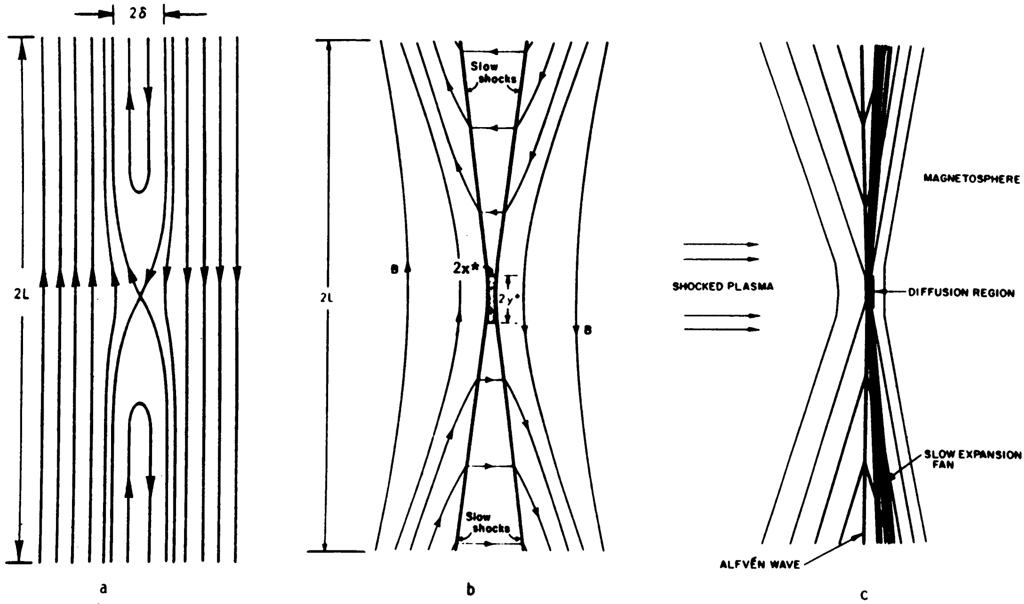

Figure 2. (a) Sweet-Parker model. Slow plasma inflow from the left and right with rapid outflow at top and bottom. (b) Petschek model. Slow plasma inflow from both sides. Acceleration into the wedge-shaped outflow regions by slow shocks. (c) Levy et al. model. Plasma inflow from the left only, with vacuum conditions on the right. Plasma acceleration into the outflow at top and bottom by large-amplitude Alfvén waves (rotational discontinuities). (After Petschek, 1964.)

outflow velocity is

$$v_{out} = B_\infty/\sqrt{\mu_o \rho} \equiv v_{A\infty} \tag{4}$$

In other words, the outflow velocity is equal to the Alfvén speed $v_{A\infty}$ based on the inflow conditions.

Second, the mass conservation law with constant density yields $v_{in}L = v_{out}\delta$ where δ and L are the half width and half length of the layer, as shown in the figure. Thus the ratio of these two lengths is

$$\delta/L = v_{in}/v_{out} = v_{in}/v_{A\infty} \equiv M_{A\infty} \tag{5}$$

where $M_{A\infty}$ is the nondimensional reconnection rate. Equation (5) indicates that the larger the reconnection rate, the thicker the layer must be (for given L).

Third, Ohm's law has the form $\underline{j} = \sigma\underline{E}$ near the separator where the plasma is semistagnant and has electrical conductivity σ. According to Ampère's law, we have $j \simeq 2B_\infty/2\mu_o\delta$ and since $E = v_{in}B_\infty$ Ohm's law yields

$$\mu_o \sigma v_{in} \delta = 1 \tag{6}$$

In other words, the magnetic Reynolds number based on the current layer half width δ and the inflow speed v_{in} is equal to unity. This statement expresses the fact that in a steady state the width δ must be such that the inflow speed is equal and opposite to the resistive diffusion speed.

If the width δ is eliminated between (5) and (6), one obtains the well-known Parker formula for the reconnection rate

$$M_{A\infty} = R_{mL}^{-\frac{1}{2}} \tag{7}$$

where R_{mL} is the magnetic Reynolds number based on $v_{A\infty}$ and L, i.e., $R_{mL} = \mu_o \sigma v_{A\infty} L$. Since the value of R_{mL} is extremely large in most cosmic applications, mostly because L is large, one is tempted to conclude from this analysis that reconnection is an insignificant process in cosmic physics. However, as discussed in the next section, cosmic plasmas may have the ability to manufacture small-scale dissipative structures, i.e., L values that are sufficiently small to permit $M_{A\infty}$ values of order unity.

The formula (7) also appears to be relevant to reconnection in tokamaks. Park et al. (1983) have generalized (7) to include the effects of viscosity, their result being

$$M_{A\infty} \sim R_{mL}^{-\frac{1}{2}}(1+\mu_o \sigma \nu)^{-\frac{1}{4}} \tag{8}$$

where ν and σ are the kinematic viscosity and the electrical conductivity, respectively. According to Spitzer (1962), the product $\mu_o \sigma_\perp \nu_\perp$, evaluated in transport perpendicular to a strong magnetic field, is

$$\mu_o \sigma_\perp \nu_\perp = \frac{3\sqrt{2}}{40\pi} (T_e/T_i)^{\frac{1}{2}} (m_i/m_e)^{\frac{1}{2}} \beta_e \tag{9}$$

where β_e is the ratio of electron pressure to magnetic pressure. This formula indicates that the modified Parker scaling (8) should be used whenever the β_e value of the plasma is of order unity or greater, as is frequently the case in cosmic plasmas.

In two-dimensional tokamak computer simulations performed by Park et al. (1983), it appears that in the m = 1 flip (internal kink) the reconnection rate is governed by (8).

A family of exact solutions of the MHD equations with constant density has been described by Sonnerup and Priest (1975) for the case of resistive stagnation point flow at a current sheet which is essentially the Sweet-Parker geometry. These authors also formulated the problem for the case where viscosity is important.

4. PETSCHEK MODEL

In order to overcome the difficulty with the small reconnection rate in the Sweet-Parker geometry, Petschek (1964) devised his now famous model in which resistive diffusion is important, not over the entire length, 2L, of the current layer, but only over a short distance, 2y*, around the separator, as illustrated in Figure 2b. In the remainder of the flow field, electromagnetic energy is converted to plasma kinetic energy and heat in a set of standing slow-mode shocks originating at the separator.

The region surrounding the separator in which resistive diffusion leads to a violation of the frozen magnetic field condition is called the "diffusion region." It has cross section 2x* by 2y* as shown in the figure. The Parker formula (7) applies to this region but with L replaced by the small length y*. The nature of the processes in the diffusion region in collision-free plasmas will be discussed in Section 6.

Away from the immediate vicinity of the separator, the frozen magnetic field condition holds except in the slow shocks. The constancy of the electric field component tangential to a shock surface guarantees that in the xy plane the magnetic field appears to be frozen across the shocks. However, the charged particles undergo a displacement in the z direction as they cross the shock whereas there is no corresponding displacement of the magnetic field lines. Thus the frozen-field condition is in fact violated in the shocks. The z displacement of the charged particles is in the direction of the reconnection electric field, $E_{\parallel}$, and it is therefore the means by which particles are energized in the shocks. Further discussion of this effect is given in Section 5. Here we simply note that the outflow speed of the plasma in the narrow wedges between the shocks can be obtained directly from the conservation laws for mass, magnetic flux, and tangential (y) momentum:

$$\left.\begin{aligned} v_{in} L &= v_{out}\delta \\ B_{in}\delta &= B_{out} L \\ \rho v_{in} v_{out} &= B_{in} B_{out}/\mu_o \end{aligned}\right\}$$

By elimination of v_{in} and B_{out} between these equation one obtains

$$v_{out} = B_{in}/\sqrt{\mu_o \rho} \tag{10}$$

As in the Sweet-Parker model, the outflow velocity is equal to the Alfvén speed based on the magnetic field in the inflow. This result is independent of the reconnection rate.

In qualitative terms, the geometrical behavior of Petschek's reconnection model for different reconnection rates is as follows. For

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very small rates the diffusion region length y* is equal to L in which case no shocks develop and the geometry is that of the Sweet-Parker model. As the rate increases, y* and x* decrease and shock pairs appear on the upper and lower sides of the diffusion region. The wedge angle between these shocks is initially very small. As the rate increases further, y* and x* continue to decrease and the wedge angle between the shocks increases. At the maximum reconnection rate, given by Petschek as

$$M_{A\infty} = \pi[8\ln(2M^2_{A\infty}R_{mL})]^{-1} \tag{11}$$

(for the incompressible case; a correction by Vasyliunas (1975) has been included), the diffusion region (y*) may be extremely small and the wedge angle substantial. The logarithmic dependence of $M_{A\infty}$ on R_{mL} permits reconnection rates $M_{A\infty} \simeq 0.1\text{-}0.2$ in typical cosmic applications.

Note that the maximum reconnection rate given by (11) is the Alfvén Mach number far upstream of the reconnection region where Petschek assumed the flow and magnetic field to be uniform. If a reference point on the x axis immediately outside the diffusion region is used instead, one finds the magnetic field to be substantially weaker there and the inflow speed substantially larger so that the inflow Alfvén Mach number is of order unity. Thus the logarithmic factor in (11) is the result of the specific upstream boundary conditions used by Petschek. Vasyliunas (1975) has pointed out that these conditions correspond to fast-mode expansion in the two inflow regions. Other boundary conditions (Sonnerup, 1970; Yeh and Axford, 1970) may lead to different upstream values $M_{A\infty}$ and different behavior (slow-mode expansion) in the inflow regions. But the result that the maximum reconnection rate corresponds to an Alfvén Mach number of order unity immediately adjacent to the reconnection region is likely to be valid regardless of the external boundary conditions on the inflow side, as long as the outflow remains unimpeded. For this reason the Petschek reconnection rate is often quoted simply as $M_A \simeq 1$.

Park et al (1983) have generalized this rate to include viscous effects, the result being

$$M_A \sim (1+\mu_o\sigma\nu)^{-\frac{1}{2}}. \tag{12}$$

This rate is not observed in computer simulations of tokamak reconnection. The reason is that the Petschek model assumes free and unimpeded outflow from the reconnection region, a condition not satisfied in tokamaks but likely to be valid in a number of cosmic situations, e.g., reconnection at the earth's magnetopause.

A precise mathematical analysis of the Petschek model may be found in Soward and Priest (1977;1982) and Soward (1982).

Petschek's model can also be developed for the case of magnetopause reconnection where the plasma and magnetic field conditions are dissimilar on the two sides of the current sheet. The special case with a vacuum on one side of the layer was discussed by Levy et al. (1964) and is shown in Figure 2c. In general, asymmetric models contain a rotational discontinuity across which most of the requisite field direction change occurs. In addition, slow shocks or slow expansion fans may be present.

Note that even a small asymmetry in the plasma density on the two inflow sides will lead to the occurrence of a rotational discontinuity. The reason for this is as follows. The slow shocks on the low density side of the configuration will have to be somewhat weaker and those on the high density side somewhat stronger than in the symmetric model in order to match flow velocities and field directions in the outflow wedges. But the strongest permissible slow shock is a switch-off shock in which the field on the downstream side of the shock is along the shock normal. If an even stronger tangential momentum change is needed on the high density side, it has to be provided by a rotational discontinuity initially followed by a slow shock which reduces the excessive momentum change provided by the rotational discontinuity, as illustrated in Figure 3.

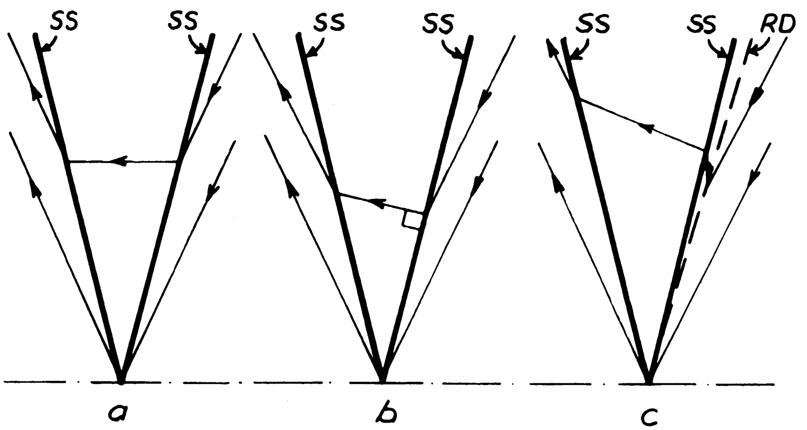

Figure 3. Asymmetric reconnection. Upper half of configuration is shown. (a) Symmetric case. (b) Density increased on the right and decreased on the left until right-hand slow shock (SS) is a switch-off shock. Left-hand shock gets weaker. (c) Density asymmetry increased further. A rotational discontinuity (RD) appears on the right, followed by a slow shock. Left-hand slow shock weakens further.

5. PARTICLE ACCELERATION IN CURRENT SHEETS

The classical reconnection models discussed in the two previous sections are magnetohydrodynamic in nature. However, the magnetospheric plasma is collision-free and it is not entirely clear how this influences the geometry or how slow shocks manifest themselves in such a medium. This point is illustrated in Figure 4 which shows symmetric

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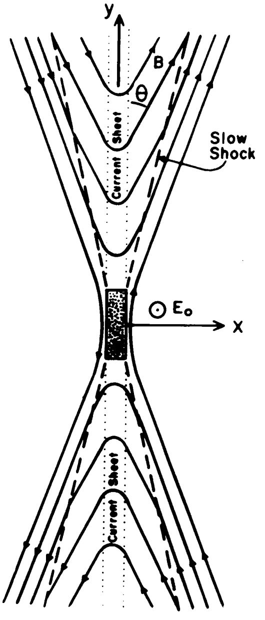

Figure 4. Hill's collision-free reconnection model. The outflow wedges are located between the central current layer and the slow shocks.

collision-free reconnection in the geomagnetic tail, as envisaged by Hill (1975). Individual particles interact with, and are energized at, a central current sheet (which in MHD terminology is a contact discontinuity and can exist only when the pressure is suitably nonisotropic), being either reflected or transmitted there. The outflowing particles form wedge-shaped regions on either side of the current sheet, regions that are somewhat similar in nature to the so-called foreshock region upstream of the earth's bow shock. In these wedges, the plasma consists of a mixture of particles that flow slowly toward the central current sheet and particles that have already interacted with it and have gained kinetic energy during the interaction. This latter population is streaming with a large field-aligned velocity component away from the diffusion region at the center of Figure 4. The outer edge of each outflow wedge is marked by particles that have passed through the diffusion region. At this location the magnetic field changes direction and magnitude in order to accommodate the higher plasma pressure in the wedge. This discontinuity is presumably a slow shock, although the usual jump conditions probably need to be modified to allow for heat flow behind it.

Since only a very small portion of the inflowing particles passes through the diffusion region, while the overwhelming majority interacts with the central current sheet, it is logical first to examine the interaction of particles with a one-dimensional laminar current sheet which may be either a shock, a rotational discontinuity or a contact discontinuity.

The first important point is that such current sheets have a non-vanishing normal magnetic field component. This feature allows one to transform away the reconnection electric field, $E_{\parallel}$, by examining the particle orbits in a frame of reference that slides along the current sheet, away from the diffusion region, with speed $v_t = E_{\parallel}/(B\sin\theta)$ where θ is defined in Figures 4 and 5. This frame is often referred to as the de Hoffmann-Teller (dHT) frame. The principal advantage of this procedure is that, in the moving frame, a particle either conserves its energy or changes it in a known manner in response to the electric potential structure, $\Phi(x)$, of the current sheet. In particular, reflected particles must have the same energy before and after reflection.

As illustrated in Figure 5 for a symmetric field-reversing current sheet with $\Phi = 0$, a particle moving with guiding-center velocity $\underline{v}_{1g}$ toward the sheet in the stationary frame appears to be moving along the magnetic field with guiding-center velocity v'_{1g} and pitch angle α_1 in the dHT frame (Fig. 5a). Assume that it traverses the sheet and leaves on the other side, moving along $\underline{B}$, with guiding-center velocity v'_{2g} and pitch angle α_2 (Fig. 5b). Conservation of energy in the dHT frame implies that $v'_{1g}/\cos\alpha_1 = v'_{2g}/\cos\alpha_2 \equiv v'$. In the stationary frame (Fig. 5c), we then see that v_{2g} can be much larger than v_{1g}. The energy change in this frame is $\Delta\varepsilon \cong m(v_2^2-v_1^2)/2$ where (for both subscripts 1 and 2) $v^2 = v_g^2 + (v'_g\tan\alpha)^2$. It is easy to show from the triangles in Figure 5 that the energy increase may be written

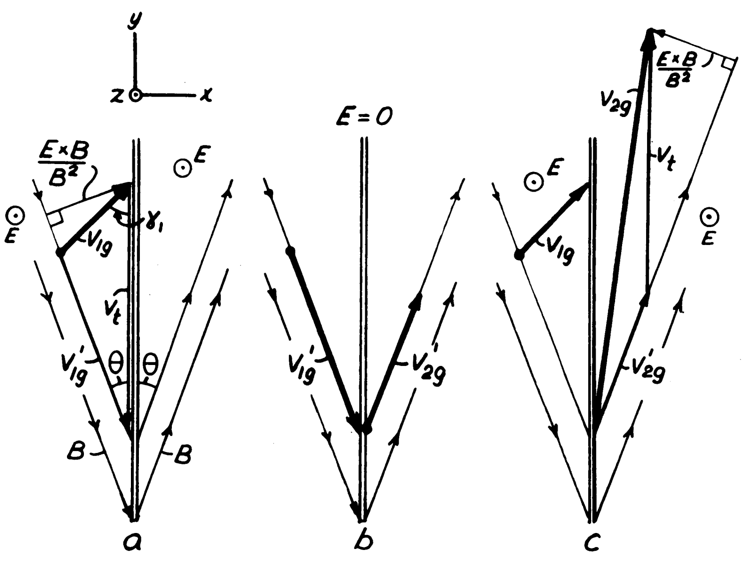

Figure 5. Particle energization in a current sheet. (a) Particle moves towards sheet with guiding center speed v_{1g}. Transformation velocity v_t along the sheet is added so that in the moving (dHT) frame the velocity is v'_{1g} along $\underline{B}$. (b) In the moving frame the electric field $\underline{E}$=0. The particle exits with velocity v'_{2g} along $\underline{B}$. (c) Return to stationary frame in which the exit velocity v_{2g} is large.

$$\Delta\varepsilon = m \frac{E_{\parallel}}{B\sin\theta} v'\cos\theta(\cos\alpha_1 - \cos\alpha_2) \quad (13)$$

or

$$\Delta\varepsilon = |q| E_{\parallel} R_{Lx} \cos\theta(\cos\alpha_1 - \cos\alpha_2) \quad (14)$$

where $R_{Lx} \equiv mv'/|q|B_x$, q being the particle charge and B_x the normal magnetic field component. Thus R_{Lx} is the particle gyroradius in the normal field.

Several general comments should be made:

(i) The particle energization is proportional to the particle mass. Thus electrons pick up very little energy and different species of ions pick up energy in proportion to their mass. At the earth's magnetopause and magnetotail where trace amounts of a variety of ion species may be present, this feature provides an opportunity for a persuasive test of the reconnection hypothesis.

(ii) The energization formulas (13,14) work for reflected as well as transmitted particles. Generalization to cases where the magnetic field and plasma flow on the two sides of the layer do not lie in a single plane is straightforward.

(iii) The energization formulas suggest large energization when θ is small. This is indeed the case for particle reflection in the earth's bow shock (e.g., Sonnerup, 1969) where the electric field is fixed. However, in the reconnection case, self consistency of magnetic fields and plasma motion implies that for a typical plasma ion $v'_{1g} = v'\cos\theta\cos\alpha_1 \simeq v_A$. For small angles θ and with the angle γ_1 in Figure 5 sufficiently different from 0 or π, we then have approximately $v_t = E_{\parallel}/(B\sin\theta) \simeq v_A$. In other words, $E_{\parallel}$ is proportional to $\sin\theta$ and the energization of a typical plasma particle is independent of θ and of the reconnection rate $E_{\parallel}$. Since $v_t \simeq v_A$ it is clear that for such a particle $v_2 \simeq 2v_A$. As noted below, these results agree with the predictions of MHD theory in which the geometrical effects discussed here are embodied in the law of conservation of tangential momentum (see e.g., Hudson, 1970; Sonnerup et al., 1981).

(iv) For the symmetric contact discontinuity in the geomagnetic tail (Fig. 4), the flow in the exit wedges contains a mixture of incoming particles with essentially zero flow speed ($v_{1g} \sim 0$) and outflowing particles with speed $v_{2g} \simeq 2v_A$ directed nearly along $\underline{B}$. This feature cannot be described by the MHD model, but the average velocity in the outflow regions is v_A in agreement with that model. For the dayside magnetopause, the magnetosheath plasma flows across a rotational discontinuity and, in the MHD model, acquires a tangential velocity of $2v_A$ in agreement with the single particle results given above. As long as only average plasma properties are considered, the MHD model gives the correct results but detailed distribution functions such as are now measured in the geomagnetic tail and at the magnetopause can of course not be obtained from that model.

(v) The single particle considerations discussed here, along with actual orbit calculations in a model tail current sheet, have been used by Lyons and Speiser (1982) to predict distribution functions in the exit wedges and good agreement with observations is obtained. To date, the corresponding calculations have not been performed at the magnetopause which usually has a much more complicated structure. However, on the fluid level, agreement between theory and observations is reasonably good (Paschmann et al., 1979; Sonnerup et al., 1981). Observations include energization of transmitted as well as reflected magnetosheath ions. Recently, the full fluid energy balance has been checked (Paschmann, private communication, 1983). It has been found that, in addition to the electromechanical energy conversion described above, the plasma enthalpy and entropy increase substantially in a rotational discontinuity. Furthermore, reflected ions may occasionally provide an important heat flow away from the discontinuity.

(vi) The energization of a particle during a single encounter with a current sheet is relatively small. In order to increase the energy to large values, multiple encounters are needed. In the geomagnetic tail, this may occur as a result of reflection near the earth in the magnetic mirrors provided by the dipole field. In other circumstances, scattering due to electromagnetic irregularities may bring some particles back to the current sheet.

(vii) One may conclude from (14) that, during its interaction with the current sheet, a particle must experience a displacement Δz in the z direction, i.e., along $E_{\parallel}$ given by

$$\Delta z = R_{Lx}\cos\theta(\cos\alpha_1 - \cos\alpha_2) \qquad (15)$$

This displacement can indeed be derived directly from the conservation of energy in the dHT frame and the conservation of the generalized particle momenta associated with the two cyclic coordinates tangential to the current sheet (Cowley, 1978). The result (15) is represented graphically in Figure 6 for the case of reflected particles. Looking along the magnetic field, one finds that the particle orbit before and after reflection must be located inside an ellipse of major and minor axes $2R_{Lx}$ and $2R_{Lx}\sin\theta$. Inside the ellipse are circles, each labeled with a specific pitch angle α, which represent the projections of the helical

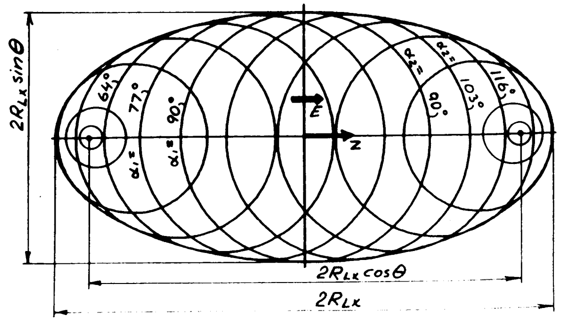

Figure 6. View along B of particle reflection in a current sheet. Incident helical orbit is a circle on the left, labeled by α_1; exit helix is circle on the right, labeled by α_2. Guiding-center displacement is purely along z (see Fig. 5).

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particle orbits. The centers of these circles, i.e., the guiding centers, are located at $z = \pm R_{Lx}\cos\theta\cos\alpha$ and, except for small pitch angles, the circles touch the ellipse. The two foci of the ellipse occur at $z = \pm R_{Lx}\cos\theta$ and correspond to pitch angles 0° and 180°. As a particle approaches the current layer, it moves on that circle, centered at $z<0$, which corresponds to its pitch angle α_1 (<90°). As it leaves the current sheet after reflection, it moves on a circle, centered at $z>0$, which corresponds to its pitch angle α_2 (>90°). Maximum displacement of the guiding center occurs for $\alpha_1 = 0^\circ$, $\alpha_2 = 1{\cdot}0^\circ$ in which case the orbit moves from the left-hand to the right-hand focus of the ellipse. Similar diagrams may be constructed for transmitted particles both in the planar case (e.g., Fig. 5) and in the case where the rotation angle of the tangential magnetic field is arbitrary. Except for thick current layers where $\alpha_1 = \alpha_2$, the actual relationship between the entry and exit pitch angles must be obtained from detailed orbit calculations.

6. DIFFUSION REGION

In discussing various physical processes in the diffusion region, it is desirable to examine the generalized Ohm's law:

$$\underline{E} + \underline{v}\times\underline{B} = \underline{j}/\sigma + (\underline{j}\times\underline{B})/ne - (\underline{\nabla}\cdot\underline{\underline{P}}_e)/ne + \{\partial\underline{j}/\partial t + \underline{\nabla}\cdot(\underline{j}\underline{v}+\underline{v}\underline{j})\}m_e/ne^2 \qquad (16)$$

Away from the separator and the slow shocks, the electric field $\underline{E}$ is balanced by $\underline{v}\times\underline{B}$, i.e., the frozen magnetic field condition holds. In the diffusion region, $\underline{E}$ remains nonzero but, since the plasma is semi-stagnant, the $\underline{v}\times\underline{B}$ term is small. Additionally, for $B_{\parallel} = 0$ the magnetic field vanishes at the separator. It is then necessary that one or more of the terms on the right-hand side of (16) becomes sufficiently large to balance $\underline{E}$. In a collisional plasma the term $\underline{j}/\sigma$ provides this balance. To do so, the current density $\underline{j}$ must be large when the conductivity σ is large; this leads to a very narrow diffusion region (small x^*), as discussed in Section 4.

In a collision-free plasma an effective resistivity $1/\sigma$ may be provided by plasma microinstabilities driven by the strong currents or gradients in the diffusion region (J.D. Huba, this Symposium). Another possibility is that electron inertial or pressure effects contained in the last two terms on the right in (16) are dominant, as discussed by Vasyliunas (1975) and Sonnerup (1979). There is no doubt that these terms must become important at the separator itself for without them, and with $\sigma = \infty$, the generalized Ohm's law implies that the magnetic field is frozen into the electron fluid which would make reconnection impossible. These terms produce the electron inertial length, λ_e, and gyroradius, R_{Le}, as important scale sizes in the diffusion region.

However, the question remains whether _all_, or even most, of the diffusion-region current is carried by electrons so that $x^* \sim \lambda_e$ (or R_{Le}).

Sonnerup (1979) has argued that the Hall term $(\underline{j}\times\underline{B})/ne$ plays an important role and leads to a much larger diffusion region width, $x^*\sim\lambda_i$ (or R_{Li}), with a substructure of size $\sim\lambda_e$ at the separator. In this situation, most of the current in the diffusion region is carried by ion drift in the z direction rather than by electron drift. Another effect is the formation of Hall-current loops in the xy plane with associated regions of positive and negative B_z values in the diffusion region: this may have a bearing on the observed magnetic field orientations in the plasmoid formed in the geomagnetic tail during reconnection (J. Birn, this Symposium). A recent analysis of the Hall effect in collisional tearing (Terasawa, 1983) is also relevant.

The inertial effects of electrons and ions discussed above can be dealt with in a qualitative manner in terms of an "inertial" conductivity $\bar{\sigma} = ne^2\tau/m$ where τ is the residence time of a particle in the diffusion region, estimated as $\tau = x^*/v_{in}$. Thus the condition $\mu_o\sigma x^* v_{in}=1$, which for constant σ leads to $x^*\sim v_{in}^{-1}$, in this case produces a value of x^* that is independent of v_{in}, i.e., of the reconnection rate. This value is

$$x^* = \sqrt{m/\mu_o ne^2} \tag{17}$$

which is the definition of the inertial length of electrons or ions depending on whether the mass m_e or m_i is used.

Microinstabilities may well be present in the diffusion region and may provide an important signature of this region and the separatrix surfaces attached to it (Scudder et al., 1983). But it is not clear that such instabilities play a significant role in allowing reconnection to occur in a collision-free plasma. The argument that in their absence electrons would be able to move quickly along the separator to cancel out $E_\parallel$ (in particular when $B_\parallel \neq 0$) fails to take account of the short residence time τ of most particles in the diffusion region. For the same reason, the diffusion region may not be a prodigious source of high energy particles, accelerated by a single large displacement along the separator. A detailed study of particle orbits near the separator would be of great interest, in particular for $B_\parallel \neq 0$, but a realistic model of the electric and magnetic fields should be used ($E_x, E_y \neq 0$; $B_z=B_z(x,y)$).

7. NONSTEADY EFFECTS

A basic discussion of nonsteady reconnection must start with an examination of the tearing mode, a topic dealt with elsewhere (J.F. Drake, this Symposium). Thus my presentation will be limited to several remarks about observed, nonsteady aspects of magnetopause reconnection.

Although the situation at the magnetopause would seem ideal for the occurrence of reconnection in a quasisteady state, whenever the interplanetary magnetic field has a substantial southward component, $B_z<0$, observations indicate that the process is mostly patchy, or at least limited to a narrow longitude segment, and that it is highly time-dependent. It is not entirely clear whether the recently discovered "flux transfer events" correspond to localized holes in the magnetopause moving away from the equatorial region, as visualized by Russell and Elphic (1978) (see also C.T. Russell, this Symposium), or to field lines connected across an open strip of the magnetopause, with the observation site located away from that strip. In either case, the observations suggest the existence of a threshold, other than $B_Z<0$, for the onset and switch off of magnetopause reconnection. The nature of this threshold is not understood but the following scenario may account for the observations.

Assume that a bundle of interplanetary magnetic field lines gets hung up, perhaps in a small indentation, over the subsolar magnetopause. The plasma will escape from the bundle by streaming along $\underline{B}$, the result being a lowering of the plasma density n, of the β value (in particular $\beta_{\parallel}$), and of the Alfvén Mach number, M_A. It may be argued that each of these factors is conducive to the onset of reconnection over the narrow longitude segment occupied by the bundle. As soon as reconnection has started, two effects occur: a deepening indentation in the magnetopause develops; and the region originally occupied by the bundle gets replenished with fresh magnetosheath flux and plasma in which n, β, and M_A return to their normal values. The latter effect may lead to the switch-off of reconnection, while the former creates a suitable site at which a new magnetosheath field bundle may get hung up. It seems possible that for varying plasma conditions, this kind of model may lead either to a succession of flux transfer events, to the occurrence of quasisteady reconnection in a narrow longitude segment, or to reconnection that ultimately spreads over a substantial longitude segment. Detailed theoretical and observational studies guided by this scenario should provide important insights into reconnection in its magnetopause version.

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DISCUSSION

Sturrock: In order to explain a solar flare, we need a system which exhibits an "explosive" or "hard" instability. What do we know about the conditions which determine whether or not reconnection is an explosive process?

Sonnerup: In my view, not much. A few suggestions for explosive instability behavior have been made. For example, explosive behavior of the ion tearing mode has been predicted by Galeev and coworkers. However, this result has been contested by Pellat.

On the whole, I believe that the type of behavior observed in solar flares is not determined solely by local conditions at the reconnection line, but by the global configuration and the dynamic accessibility of the free energy stored in it.

Mahajan: What is the definition of a "classical picture"? Does it have anything to do with the form of the Ohm's law?

Sonnerup: No, it does not refer to classical versus anomalous resistivity. I used the phrase "classical reconnection models" for the Sweet-Parker and the Petschek models.

Vasyliunas: On the definition of reconnection: 1) As you pointed out, definitions in terms of plasma flow or in terms of electric field are practically equivalent; the advantage of the former is that the trivial vacuum case is excluded. 2) It may be preferable to define reconnection globally, as palsma flow across an electric field E in the separatrix surface; the existence of E along the separator line can then be deduced as a consequence. 3) Various terms: reconnection, field line merging (preferable in my opinion, for linguistic reasons), field annihilation, etc. are all synonymous.

Sonnerup: (1) I prefer to think of flux transfer as the principal characteristic of reconnection. Since such transfer occurs unimpeded in a vacuum, that case must of course be included in the definition; I do not see that as a disadvantage. The definition in terms of plasma flow across a separatrix does, on the other hand, have the disadvantage that the case of annihilation of exactly antiparallel fields (e.g., the stagnation point flows studied by Priest and myself a few years ago) does not qualify as reconnection under such a definition. (2) I actually prefer the local definition because the reconnection actually takes place at the separator. From this definition one can then deduce (with minimal additional assumptions) that an electric field is also present elsewhere on the separatrix surface. (3) It is true that the terms reconnection, merging, and annihilation are used more or less interchangeably. However, in my view it would be desirable to make a slight distinction: reconnection describes the case of a distinct separator line; annihilation describes the case where the separator has degenerated to a surface (this occurs in a current sheet without a normal magnetic field component). Finally, merging could be used to incorporate both of the preceding situations.

Bratenahl: I am fascinated with Nancy Crooker's idea that the maximum reconnection rate occurs when the fields are antiparallel. With the new calculational machinery, are we on the track of being able to settle this appealing idea of Crooker?

Sonnerup: Ordinary reconnection theory predicts a simple formula for the reconnection electric field as a function of the angle between the two reconnecting magnetic fields. This formula gives a maximum when the fields are anti-parallel. Nancy Crooker's proposal is more radical. She argues that reconnection may occur only when the two reconnecting magnetic fields are exactly, or very nearly, antiparallel. I do not know of any strong theoretical or observational support for this idea. In a tokamak the reconnecting fields form only a very small angle. On the other hand, the recent collision free electron tearing mode analysis by Coroniti and Quest does indicate a fairly strong dependence of the growth rate on the angle between the fields.

A definite answer to the question must await a better understanding of the role played by B in the diffusion region.

Van Hoven: Can the width of diffusion region be uniquely specified for the case of steady reconnection?

Sonnerup: In the case where classical resisitivity η multiplied by the current density is used to balance the electric field in the diffusion region, the width x^* is such that the magnetic Reynolds number $M_o v_1 x^*/\eta \simeq 1$. In that case x^* simply becomes as small as is required

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in order to maintain this approximate equality. If, on the other hand, the effective resitivity is a function of current density or of the reconnection rate, then the situation may be different. For example, in my paper I have shown that the width of a diffusion region in which ion inertial effects dominate is always of the order of the ion inertial length λ; (or perhaps the ion gyroradius). Similarly, a diffusion region dominated by electron inertial effects would have width $x^* \simeq \lambda_e$.

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