THEORETICAL APPROACHES TO THE DESCRIPTION OF MAGNETIC MERGING: THE NEED FOR FINITE β_e, ANISOTROPIC, AMBIPOLAR HALL MHD

J. D. SCUDDER
Department of Physics and Astronomy, University of Iowa, Iowa City, USA

Received April 28, 1997; Accepted in final form May 16, 1997

Abstract. Theoretical models of magnetic reconnection are reviewed with a critical view of their suitability for astrophysical plasmas, with a focus on those sampled plasmas near the magnetopause. Frequently the approximations are more those of convenience than physically justified. It is concluded that magnetic reconnection cannot be qualitatively or quantitatively addressed with any one fluid MHD picture unless the Hall, ambipolar and inertial emfs are included in the Generalized Ohm's Law. The observed size of electron pressure anisotropies ensures that the thawing of magnetic flux is almost always determined by the often neglected ambipolar term of the Generalized Ohm's Law. Thus resistive MHD or even resistive Hall MHD cannot possibly give a correct structural picture of the reconnection current carrying layer at the magnetopause. In the magnetotail the ion inertial "resistivity" is much larger than coulomb resistivity with a similar structural form as the coulomb emf. However, until recently the ambipolar contributions there have not been considered. This change in viewpoint of the controlling factors for thawing of magnetic flux parallels the recent evolution of understanding of collisionless shocks, where initially stochastic wave-resistivities were thought to substitute for the coulomb dissipation of high density shock waves. Now these shocks are known to be controlled by coherent agents that can modify emf's such as the ambipolar electric field, the Hall contributions of the gyrating ions, and the electric electron drift in the shock layer to support the current without thawing flux and without any requirement of ohmic dissipation *per se*. The observational tests that reconnection is a viable process for plasma entry in the magnetosphere are briefly reviewed. Sites where these conservation laws are said to be approximately fulfilled are discussed with an eye toward systematic experimental issues of these tests. That magnetic shear poorly indexes "good" Walén testing layers may be an indication that the resistive dissipation is either not uniformly important across the data set or resistive emf's are not the appropriate agent for the thawing of flux. The ambipolar scale length clearly exceeds the resistive or electron skin depth regime with layers that pass the "good" Walén test layers which have $\beta < 5$; this may indicate the importance of the ambipolar violations to the frozen field description.

1. Introduction

Based on circumstantial evidence the space physics community is convinced that **B** field line topologies change about the Earth. Were this not the case the solar wind magnetic field hung up over the magnetopause would intensify indefinitely. The substorm activity on the ground and witnessed in space are indicators of the reality of this scenario. Nevertheless, we do not "know" what process or processes mediate the topology changes whereby flux is shed from the dayside magnetosphere and recycled to the lobes in the magnetotail. In large scale plasmas these changes in topology are accompanied by specific rearrangements of the plasma. The effects of field line interconnection that accompany these topology changes have been modeled most simply with fluid-like equations of MHD by adding lumped parameters

Space Science Reviews **80**: 235–267, 1997.

such as localized (anomalous) enhancements of resistivity. While this approach may be the simplest, there remains the problem of defining the agent for this resistivity. There are, however, many other *possible* ways that the reconnection can be permitted that this paper will summarize. At least two of the known alternative agents for enabling reconnection are present even in the collisionless limit and have a wholly different structure than that mocked up by a collisional or enhanced *ad hoc* resistivity.

A number of excellent reviews on the description of merging already exist from the fluid picture (cf. Axford 1984; Vasyliunas, 1975; Sonnerup, 1984; Lee, 1995), and the collisionless picture (Hill, 1975; Cowley, 1982), so a brief statement of purpose for the present paper is appropriate. The present situation in the reconnection arena is much like the situation a few years ago with collisionless shocks. Then, observers told theoreticians shocks were present; theoreticians gave an explanation in terms of anomalous resistivity driven by micro-instabilities in the collisionless shock current layers. In this way a one for one substitution between two body resistivity and collective resistivity yielded a theory to "explain" the existence of the shock with insufficient empirical corroboration for the assumptions of their theory. However, with further experimental scrutiny it became clear, contrary to this *explanation*, that a rather diverse set of neglected *coherent* effects in the shock layer could explain their existence and the reformation of the shock layers themselves without this substitution of waves for binary collisions. Further, careful determinations of scales and sources of free energy found that the picture of linear growth, followed by non-linear reaction thought to facilitate anomalous resistivity, was a cycle that had insufficient time to take place before the fluid was swept through the layer and no longer able to tap the free energy. In the present view of shocks the collective wave particle effects "round the corners" of otherwise coherent processes in the DC electromagnetic fields of the shock structure. The self-consistent, post MHD, electromagnetic field including $E_{||}$ was crucial to that paradigm shift.

The physics of reconnection is almost totally determined theoretically by an acceptably accurate description for the parallel electric field. Resistivity is just one of several ways to support parallel electric fields. Because reconnection involves moving plasmas in the presence of magnetic fields at current sheets with gradients, this parallel electric field will certainly have a coherent ambipolar component; whether a sufficiently vigorous stochastic $E_{||}$ can be found to substitute for the almost negligible classical coulomb resistivity and dominate the coherent ambipolar parallel electric field (which has usually been neglected) remains to be seen in the problem of astrophysical reconnection. It is clear for $\beta \simeq 1$ plasmas that there are significant coherent sources of the parallel electric field that may routinely be more important in understanding the factors that control the onset of reconnection than surmised anomalous resistivities from yet to be defined wave particle instabilities.

This review will attempt to organize all the effects that could possibly be involved in describing magnetic reconnection at the level of the conservation laws of magneto-fluid dynamics. It will become clear that demonstrably large effects have been postulated to be unimportant in much of the modeling of reconnection that has been completed to date. The main impediments are (1) the equation of state problems and (2) moment closure for an astrophysical plasma.

Magnetic reconnection is a process that may occur when the magnetized plasma is *non-ideal*. As with most adjectives, there are various levels of *non-idealness*. Broadly speaking reconnection involves loss of the line or flux preserving character of the plasma flow. Colloquially, the field line is said to be "broken" or have its extremities reconnected with a different topology. Ultimately the mixing of plasma populations from different topological field lines is the result (Vasyliunas, 1972, 1975). In this parlance reconnection occurs when plasma flows across the separatrix that bounds the topologically different regimes at the site of reconnection. Vasyliunas, Axford and others have emphasized that it is not enough for there to be local departures from ideal MHD, without simultaneous global requirements outside such "diffusive" regions that determine whether any reconnection proceeds in a quasi-steady manner. In the context of Earth there are three topologically distinct groups of field lines at any give moment: those with both foot points anchored in the sun; those with one end anchored in the sun with the other in the core of the earth; and those with both foot points anchored in the core of the earth. The time sequence of reorganization of the plasma that is first on the first topological bundle of field lines, then the second and subsequently the third requires the relaxation of the concept that field lines have an integrity throughout their dynamical history. The concept of field lines or frozen flux being preserved *within* each topological type, requires some relaxation of the field line concept in the locations where these topological changes are mediated.

These rather difficult concepts are often *summarized (or hidden?)* by saying that reconnection occurs or is possible in regimes where the "frozen field approximation"

$$\mathbf{E} = -\frac{\mathbf{U} \times \mathbf{B}}{\mathbf{c}}$$

is violated (here $\mathbf{U}$ is the center of mass velocity); or, whenever there exists $E_{||}(\mathbf{x})$ collocated with field aligned currents. Still others say that reconnection occurs when field energy is transferred to plasma energy when crossing such topological boundaries. In this state of affairs, the literature is full of rather diverse statements that approach almost religious fervor concerning whether reconnection is or is not an established operative mechanism at the earth's near environment. To witness flow across a change in topology is almost impossible with a single spacecraft. Indirect measures of such a flow, such as determining the tangential electric field at the magnetopause (Mozer et al., 1978), and the normal component of $\mathbf{B}$ (Sonnerup and Ledley, 1979), or its consequences via the Walén (1944) test (Paschmann et al. 1979, Sonnerup et al., 1981, Aggson et al., 1983; 1984) or time of flight dispersion are used as proxies for the process under models of the implementation of the interconnection. If these models are themselves flawed or over simplified, the tests of these models remain inconclusive, or, at least parametric in the premises of the model. The purpose of this review is to elucidate the governing equations of magnetic reconnection with as few simplifications as possible and then discuss the tests that have been attempted and their results in light of simplifications that have made these tests possible.

2. Frozen Flux and its Thawing

To understand the various theoretical approaches to the description of magnetic merging, or reconnection, there must first be a clear understanding of the concept of "magnetic flux" and "magnetic line" preservation. The concept of flux preservation gives rise to the mnemonic of spaghetti-like flux tubes as entities with an integrity that are locally convected by the components of the plasma flow field that are locally orthogonal to the flux tube. In this way Alfvén's (1942) concept of moving flux tubes has gained acceptance in the *cartoon* level description of the time evolving dynamics of plasma magnetohydrodynamics. The huge success of Alfvén's point of view belies the subtlety that it is only approximately true that "magnetic flux" and "magnetic line" preservation adequately describe nature. It is precisely where this approximation breaks down that magnetic merging with its capacity to change the topology of these spaghetti-like tubes is required.

Confusingly, there are yet further simplifications of the description of real plasma dynamics that compromise the idealization of flux tubes frozen into the perpendicular motions of the fluid. There are many different representations of the fluid flow fields, $\mathbf{W}$, within which the magnetic field can be viewed as approximately having a flux or field line preserving character. For the magnetic flux to be "frozen" to the perpendicular motions of these different fluid fields requires stronger or lesser ancillary *assumptions* about the plasma's ability to react to the deformations of the topologically preserved bundles of spaghetti and the plasmas found on them.

The most commonly used fluid flow field $\mathbf{W}$ is that of the center of mass, $\mathbf{U}_{cm}$, of the magnetohydrodynamic fluid defined by the relation

$$\mathbf{U}_{cm} = \frac{\Sigma_i n_i m_i \mathbf{U}_i + n_e m_e \mathbf{U}_e}{\Sigma_i n_i m_i + n_e m_e} \tag{1}$$

with the electric field determined from the **approximate** relationship of the form

$$\mathbf{E} + \frac{\mathbf{U}_{cm} \times \mathbf{B}}{c} \simeq 0 \tag{2}$$

As is well known (2) represents but a portion of the difference between the ion and electron momentum equations (cf 14a, below). Among the neglected terms are those associated with the magnetic back emf associated with currents that may be caused by the flow field $\mathbf{U}_{cm}$, any resistive emfs, any ambipolar emfs, and finite larmor radius emfs, or any emfs associated with inertia and the acceleration of the plasma. When the flow field associated with the center of mass drives spaghetti-like tubes towards one another, there are corrections to the zero on the rhs of (2) that gnaw away at the integrity of individual tubes of force and require a broader overview of the magnetic field than as a strict ensemble of spaghetti strands: the "frozen" field approximation must be modified in those regions.

As an example, the magnetopause Chapman-Ferraro current layer is a site where there are surface currents, where magnetic tubes of forces are thrust towards one another, where terms missing from (2) are required by the large scale dynamics (boundary conditions) of the solar wind impinging on the earth's magnetic field. There is ample circumstantial evidence that changes in external boundary conditions can facilitate topological rearrangement of magnetic flux inside the preexisting magnetospheric cavity. Sometimes steady and at other times patchy or time dependent, frozen flux violations implicit in magnetic reconnection are frequently invoked to explain the shedding of flux from noon to the lobes and the dipolarization of the tail magnetic field and other substorm morphology. These "explanations" are connected with the view that there are places on the Chapman-Ferraro layer where the approximation of (2) is inadequate.

Quite generally when (2) *equals* zero one can prove that magnetic field lines are preserved throughout their dynamical advection with $\mathbf{U}_{\perp,cm}$. Accordingly the topology of these lines of force, however contorted they become, cannot be modified. It is however difficult to imagine non-trivial velocity fields which do not produce distributions of current, acceleration, and pressure so that the neglected terms of (2) will become important and critical at places in the dynamics. This intrinsic incompleteness of (2) vacates any realistic expectation that tubes of magnetic force in real magneto plasma dynamical problems can have a global and time independent integrity.

To see that there may be places and scales where the concept of magnetic field or flux lines carried by the plasma does not make local sense we now review the machinery to describe the total time derivative of magnetic flux in the manner of Rossi and Olbert (1970) as amplified by Siscoe (1983). Consider as indicated in Figure 1 the magnetic flux Φ piercing the closed loop C, with local normal $\hat{\mathbf{n}}_1$ defined as:

$$\Phi = \int_{C(t)} \mathbf{B} \cdot \hat{\mathbf{n}}_1 da \tag{3}$$

The time rate of change of this flux *for a linked set of fluid elements moving with the velocity field* $\mathbf{W}$ which transports $C(t) \rightarrow C'(t + \Delta t)$ is defined to be

$$\frac{d\Phi}{dt} \equiv \lim_{\Delta t \to 0} \frac{\Phi(t + \Delta t) - \Phi(t)}{\Delta t} \tag{4}$$

As illustrated in Figure 1, the enclosed area of the loop $C(t)$ is denoted by the oriented surface labeled S_1 whose normal is on the same side of the loop C as is the local flow vector $\mathbf{W}(\mathbf{x}, t)$. The velocity field $\mathbf{W}$ displaces the original loop $C(t)$ to a new loop $C'(t + \Delta t)$ with an enclosed surface element S_2 and surface normal $\hat{\mathbf{n}}_2$

located on the same side of S_2 as $\hat{\mathbf{n}}_1$ is of S_1 relative to $\mathbf{W}$. During Δt the original loop sweeps out a new open surface which is the union of S_2 and S_3 which is also still bounded by the original curve $C(t)$. Thus the evolution of the enclosed flux after *linked advection* by the flow field $\mathbf{W}$ is

$$\frac{d\Phi}{dt} \equiv \lim_{\Delta t \to 0} \frac{\Phi_2(t+\Delta t) - \Phi_1(t)}{\Delta t} \tag{5}$$

Consider the entire closed surface of Figure 1 made up of S_1, S_2, and S_3. Gauss's law and the source free Maxwell equation implies that the sum of the magnetic fluxes projected along the *outward normal* from this closed surface should vanish at any time. Since Φ_1 is defined for an ingoing normal to the closed surface, the $\nabla \cdot \mathbf{B} = 0$ equation implies that

$$\Phi_2(t+\Delta t) = \Phi_1(t+\Delta t) - \Phi_3(t+\Delta t) \tag{6}$$

Using (6) into (5) and (3) for the remaining fluxes eliminates $\Phi_2(t+\Delta t)$ to yield

$$\frac{d\Phi}{dt} = \lim_{\Delta t \to 0} \frac{1}{\Delta t} \left[\int_{S_1} [\mathbf{B}(t+\Delta t) - \mathbf{B}(t)] \cdot \hat{\mathbf{n}}_1 da - \int_{S_3} \mathbf{B}(t+\Delta t) \cdot \hat{\mathbf{n}}_3 da' \right] \tag{7}$$

where $\hat{\mathbf{n}}_3 da'$ is the outward directed differential area normal to surface S_3 implied by the expression

$$\hat{\mathbf{n}}_3 da' = d\mathbf{C}(\mathbf{t}) \times \mathbf{W}(\mathbf{x}, \mathbf{t}) \Delta t \tag{8}$$

where $d\mathbf{C}(t)$ is the local tangent differential along C(t).

Using (8) into the last term of (7) produces the simplification

$$\begin{aligned} \int_{S_3} \mathbf{B}(t+\Delta t) \cdot \hat{\mathbf{n}}_3 da &= \int_{C(t)} \mathbf{B}(t+\Delta t) \cdot d\mathbf{C}(t) \times \mathbf{W}(\mathbf{x}, t) \Delta t \\ &= \int_{C(t)} \mathbf{W} \times \mathbf{B}(t+\Delta t) \cdot d\mathbf{C}(t) \Delta t \end{aligned} \tag{9}$$

Inserting (9) into (7) and canceling the differentials where possible yields

$$\frac{d\Phi}{dt} = \lim_{\Delta t \to 0} \left[\int_{S_1} \frac{\mathbf{B}(t+\Delta t) - \mathbf{B}(t)}{\Delta t} \cdot \hat{\mathbf{n}}_1 da - \int_{C(t)} \mathbf{W}(\mathbf{x}, t) \times \mathbf{B}(t+\Delta t) \cdot d\mathbf{C}(t) \right] \tag{10a}$$

Proceeding to the limit yields

$$\frac{d\Phi}{dt} = \int_{S_1} \frac{\partial \mathbf{B}}{\partial t} \cdot \hat{\mathbf{n}}_1 da - \int_{C(t)} \mathbf{W}(\mathbf{x}, t) \times \mathbf{B}(t) \cdot d\mathbf{C}(t) \tag{10b}$$

Using Stoke's theorem the last term of (10a) becomes

$$\int_{S_1} \nabla \times (\mathbf{W} \times \mathbf{B}(t)) \cdot \hat{\mathbf{n}}_1 da \tag{11}$$

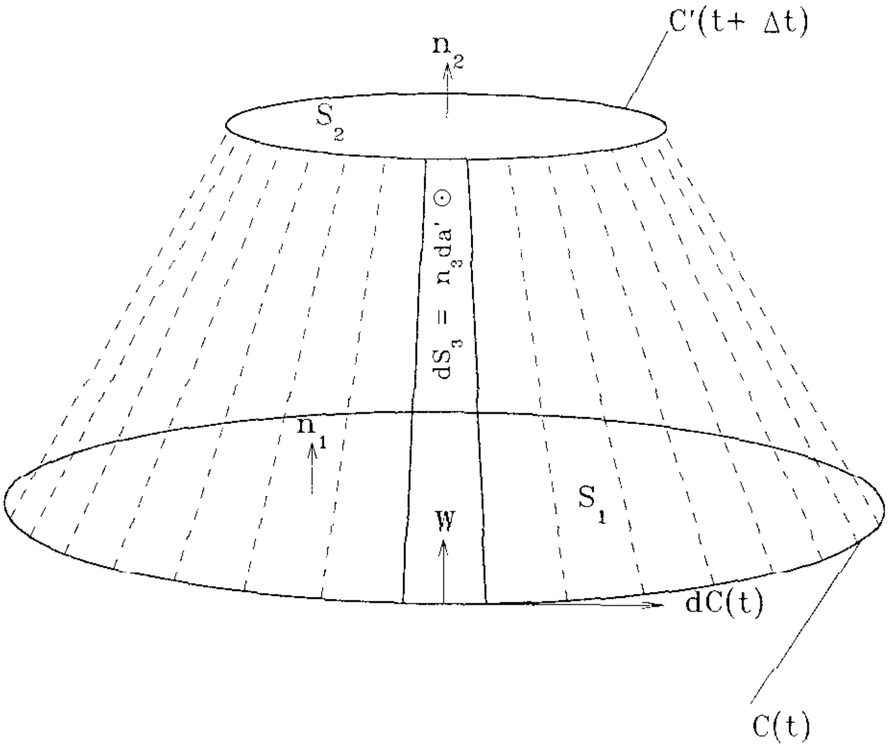

Figure 1. Representation after Rossi and Olbert, (1970) of the linked advection that maps the closed contour $C(\mathrm{x}, t)$ to a new curve $C'(\mathrm{x}', t + \delta t)$ as implemented by the flow field $\mathrm{W}(\mathrm{x}, t)$. Note that the union of surfaces S_1, S_2, S_3 is a closed surface, while S_1 and the union of S_2 and S_3 form open surfaces with C as their common bound. The surface S_3 is swept out in time by the advection of W.

while Faraday's law restructures the first term to yield the time variation of advectively connected flux:

$$\frac{d\Phi}{dt}|_{C_1} = -\int_{S_1} \nabla \times [c\mathbf{E} + \mathbf{W} \times \mathbf{B}] \cdot \hat{\mathbf{n}}_1 da \tag{12}$$

Two distinctions are usually drawn from (12) to differentiate two classes of flow fields both of which preserve flux with $\frac{d\Phi}{dt} \equiv 0$: (i) those flow fields $\mathbf{W}'$ that are flux and, more restrictively, *field line preserving* where:

$$c\mathbf{E} + \mathbf{W}' \times \mathbf{B} = 0 \tag{13a}$$

and (ii) those $\mathbf{W}''$ that are only *flux preserving*:

$$\nabla \times [c\mathbf{E} + \mathbf{W}'' \times \mathbf{B}] = 0 \tag{13b}$$

From (13a) it is now clear that the *assumption* of (2) as being true would imply that nature can only be magnetic line preserving and that magnetic field lines are carried about with any flow field whose motions perpendicular to the magnetic field are those of the fluid's center of mass:

$$\mathbf{W}_{\perp} = \mathbf{U}_{cm_{\perp}}$$

Alternatively there are a whole class of flow fields from (13b) with $\mathbf{W}''$ that while not necessarily preserving individual field lines in the same fluid element under dynamics have a balanced gain and loss of (previously labeled) field lines threading the volume in such a way that the net *number* of counted field lines, or flux, in the evolving fluid element is preserved.

3. The General Relationship for E

The magneto-fluid dynamical equations are incomplete without the relationship that is now standardly called the *Generalized Ohm's Law* which relates the electric field in the plasma to other field and plasma quantities. In a very real sense this equation is the synthesis within the one fluid picture that electrons have a different equation of motion than do the ions (cf. Rossi and Olbert, 1970).

The first and the most commonly encountered form is related to (2) and found by subtracting the electron momentum equation from that of the ions and solving for the electric field (while neglecting $\frac{m_e}{M_i}$ and space charge) with an *approximate* form (Rossi and Olbert, 1970)

$$\mathbf{E} + \frac{\mathbf{U}_{cm} \times \mathbf{B}}{c} \simeq \frac{\mathbf{J} \times \mathbf{B}}{ecn_e} - \frac{\nabla \cdot \mathbf{P}_e}{en_e} - \frac{m_e}{M_i}\frac{\nabla \cdot \mathbf{P}_i}{en_i}$$

$$+\frac{m_e}{ne^2}\left[\frac{\partial \mathbf{J}}{\partial t} + \nabla \cdot (\mathbf{U}_{cm}\mathbf{J} + \mathbf{J}\mathbf{U}_{cm})\right] + \eta \cdot \mathbf{J} \qquad (14a)$$

Alternatively, from the electron momentum equation the relationship for $\mathbf{E}$ is found to be:

$$\mathbf{E} + \frac{\mathbf{U}_e \times \mathbf{B}}{c} = -\frac{\nabla \cdot \mathbf{P}_e}{en_e} +$$

$$-\frac{m_e}{en_e}\left[\frac{\partial n_e \mathbf{U}_e}{\partial t} + \nabla \cdot (n_e \mathbf{U}_e \mathbf{U}_e)\right] - \frac{GM_* m_e \hat{\mathbf{r}}}{er^2} - \frac{m_e}{en_e} < \frac{\delta f_e}{\partial t}|_{ei} \mathbf{v}_e >_{\mathbf{U}_e} \qquad (14b)$$

The gravitational attraction with the nearest attracting object of mass, M_* at distance r to the present location is included in (14a-c). The role of gravity is implicit in (14a) through the center of mass equation of motion, while being explicitly present in the electron and ion specific forms. From the ion momentum equation comes the relationship

$$\mathbf{E} + \frac{\mathbf{U}_i \times \mathbf{B}}{c} = +\frac{\nabla \cdot \mathbf{P}_i}{en_i}$$

$$+\frac{m_i}{en_i}\left[\frac{\partial n_i \mathbf{U}_i}{\partial t} + \nabla \cdot (n_i \mathbf{U}_i \mathbf{U}_i)\right] + \frac{GM_* M_i \hat{\mathbf{r}}}{er^2} - \frac{M_i}{en_i} < \frac{\delta f_i}{\partial t}|_{ie} \mathbf{v}_i >_{\mathbf{U}_i} \qquad (14c)$$

Although not entirely equivalent (14a c) together with (12) above contribute new insight into the possible integrity of a description of magnetic reconnection that is possible with different approximations. As can be seen explicitly in (14a) (2) implies *a priori* that there are no important currents, pressure gradients or accelerations in order that the magnetic tubes of force will be adequately described as convecting with the perpendicular portions of the $\mathbf{U}_{cm}$.

Equation (14b) and (12) suggest that magnetic field *lines* are approximately frozen in the electron's center of mass frame moving with local speed $\mathbf{U}_e$ *even in the presence of currents*, provided the electron pressure gradient terms and the inertial terms are small. Equation (14b) is clearly a more general description than (2) and clearly indicates that the neglect of the Lorentz emf or Hall term in (2), or the center of mass form (14a) preempts the slippage of the electron fluid with respect to the ions (i.e. currents) from being naturally described. Thus the neglect of the Hall emf preempts important current layer phenomena. Long before physical "thawing" of flux is involved, it must be clear that the assumption of (2) to explain naturally occurring plasmas will have serious observational contradictions. Even beyond the Hall emf issue it is also clear that assuming the electrons carry the magnetic field in a frozen way ($RHS(14b) \simeq 0$) is *not* a sufficiently general point of view to compare with nature.

In principal (14c) could be used to constrain $\mathbf{E}$ as described from the ion velocity field $\mathbf{U}_i$; unfortunately (14c) describes the motion of field lines carried by the electrons using almost all the remaining ion terms including the ion inertial terms, and the ion pressure divergence; these terms are all sizeable in current layers that form as tubes of force become entangled while moving as frozen flux bundles. These current layers are usually ion gyroradius in scale in plasmas thus requiring at least Finite Larmor Radius (FLR) descriptions for the ion pressure tensor. *All* of these substantial inertial and FLR corrections are required (and hence not optional) to describe successfully the content of the generalized Ohm's law in the ion flow field's frame of reference; these terms are required with great fidelity before (14c) can emulate with precision the succinct description of the field lines moving laminarly as convected by the electrons as a group. *Most disconcerting is the realization that selectively picking some of the terms on the rhs of (14a,c) while leaving others may not even adequately describe the lowest order freezing of magnetic flux into the electron frame of reference!*

Gravity in (14b,c) also plays a role in astrophysics in the Generalized Ohm's law. Note, however, that it does not appear in (14a) (cf. Rossi and Olbert 1970) being hidden in the acceleration terms of the center of mass in (14a). Gravitational acceleration is a good example of a contributor that erodes field *line* preservation but not magnetic *flux* conservation, since the curl of the conservative force's acceleration vanishes. Thus in any of the specific fluid fields of (14b,c) identified above the role of gravity is to destroy field line preservation in the fluid frame, but not flux preservation.

The next level of sophistication in the retention of additional effects in the Generalized Ohm's Law crosses over into the regime where energy equations and closure equations for the pressures are required. Aside from the temperature dependence of the resistivity, this is the first place in the Generalized Ohm's Law that the non-zero β of the particles and their attendant acceleration from pressure gradients are factored into the electric field considerations. From one point of view the theorist or simulator can view these thermal terms as a vexing complication, to be avoided if at all possible as in Lottermoser and Scholer (1997). These terms require one or more energy equations (as in Ma and Battacharjee, 1996) which require further focus on the perplexing issues of plasma fluid dynamical closure relations. However, by observations astrophysical plasmas have significant pressures in the particles making their neglect a serious compromise for a realistic description of magnetic reconnection phenomena. Heat conduction and some forms of magnetic viscosity are transport phenomena that are almost always omitted from the discussion of the reconnection layers.

To include the so-called *ambipolar electric field* from the divergence of the electron pressure (which occurs in all three expressions above) and the ion pressure tensor that occurs in the center of mass and ion form of (14a,c) requires knowledge of the equation(s) of state in a plasma. When these plasmas are dilute as at the magnetopause there is no known equation of state that relates the electron pressure tensor to the lower order fluid moments and the magnetic field. Such difficulties do **not**, however, preempt the importance of such terms in the predictions for magnetic flux thawing.

Magnetic flux is frozen in the electron's frame of reference unless

$$\nabla \times (c\mathbf{E} + \mathbf{U}_e \times \mathbf{B}) \neq 0$$

Temporarily neglecting inertial terms, magnetic flux is thawed in the electron fluid frame whenever

$$\nabla \times \frac{\nabla \cdot \mathbf{P}_e}{en_e} \neq 0 \qquad (15a)$$

Expanding the curl reveals that this condition is equal to

$$(\nabla \cdot \mathbf{P}_e \times \nabla \ln n_e + \nabla \times \nabla \cdot \mathbf{P}_e) \neq 0 \qquad (15b)$$

In general the second term of (15b) vanishes for any isotropic electron pressure. Accordingly, the total curl of (15a) vanishes under the *further* assumption beyond isotropy that the electron plasma is *barotropic*, with its isobars and isodensity surfaces coincident. Obviously barotropic equations of state are not the most general class; the complimentary class of *baroclinic* closure equations, which do not meet the barotropic definition, are already important agents for the thawing of magnetic flux. For baroclinic isotropic electron pressure (15a) becomes

$$\nabla \times \frac{-\nabla \cdot \mathbf{P}_e}{en_e}\Big|_{baroclinic-isotropic} = \frac{1}{en_e}\nabla \ln n_e \times \nabla P_e \qquad (15c)$$

Still more complicated are *anisotropic* pressure tensors appropriate for circumstances where the electrons remain magnetized and *non-gyrotropic* equations of state when they are not. In the magnetized, anisotropic circumstance both terms of (15b) survive and may be written as

$$-\nabla \times \frac{\nabla \cdot \mathbf{P}_e}{en_e}|_{anisotropic} = \frac{1}{en_e}\left[\nabla \ln n_e \times \nabla \cdot \mathbf{P}_e - \nabla\left(\frac{\partial}{\partial s}\left(\frac{P_\parallel - P_\perp}{B}\right)\right) \times \mathbf{B}\right]$$

$$+\frac{1}{en_e}\left[\frac{\partial}{\partial s}\left(\frac{P_\parallel - P_\perp}{B}\right)\nabla \times \mathbf{B} + \nabla\left(P_\parallel - P_\perp\right) \times \frac{\partial \hat{\mathbf{b}}}{\partial s} + \left(P_\parallel - P_\perp\right)\nabla \times \frac{\partial \hat{\mathbf{b}}}{\partial s}\right] \tag{15d}$$

where

$$\nabla \cdot \mathbf{P}_e = \nabla P_\perp - \left(\left(P_\parallel - P_\perp\right)\frac{\partial \ln B}{\partial s}\right)\hat{\mathbf{b}} + \frac{\partial\left(P_\parallel - P_\perp\right)\hat{\mathbf{b}}}{\partial s} \tag{15e}$$

has been used. If the electrons are no longer magnetized then the electron pressure tensor retains no symmetries or simplifications and k'th vector component of 15a become

$$\left(\nabla \times \frac{-\nabla \cdot \mathbf{P}_e}{en_e}|_{non-gyro}\right)_k = \frac{\epsilon_{ijk}}{en_e}\left[\frac{\partial \ln n_e}{\partial x_i}\frac{\partial P_{lj}}{\partial x_l - \frac{\partial^2 P_{lj}}{\partial x_i \partial x_l}}\right] \tag{15f}$$

with Einstein summation convention implied and ϵ_{ijk} the Levi-Civita fully anti-symmetric tensor.

The general barotropic form for $\mathbf{P}_e \equiv P_e$ occurs when P_e is *solely* a function of density alone. A polytrope electron equation of state is *a* form that meets the barotropic condition. There is no known theorem that even guarantees that a plasma or even the electron part of the plasma must emulate polytrope behavior or even be a local function relating moments of the plasma in the *same* location. Polytropes have been predicted (Scudder and Olbert, 1979) and have been found to organize solar wind electrons (Sittler and Scudder, 1980). Unfortunately, there are growing indications that the plasma pressure equations for a rarefied astrophysical plasma are not universal, but context dependent (Scudder, 1992; Osherovich et al., 1993, 1995; and Fainberg et al., 1996). The documentation of polytropic behavior in the freely expanding solar wind is no *a priori* guarantee for its relevance as the equation of state within the magnetopause current layers. Such barotropic corrections to the electric field do not contribute to the erosion of magnetic *flux* but do introduce additional freedom for moving flux that otherwise would pile up; these terms also add new scale lengths (cf section 6) and time scales to the problem of reconnection (Vasyliunas, 1975; Ma and Battacharjee, 1996; and Kleva et al., 1995).

The polytropic barotropic approximation for electrons compromises the magnetic line preservation but permits magnetic flux and its topology to be preserved.

In this case the precise same field lines are not always penetrating the curve C of Figure 1 as it is carried with the electrons, but the net coming and going of lines of force and their attached electrons by pressure gradient drifts are such that the enclosed magnetic flux is invariant. Gradient and curvature drifts are compensated with isotropic pressure. An alternative viewpoint is that the magnetic field line is "carried" at the composite speed given by

$$\mathbf{W_B} = \mathbf{U}_{\perp,e} - c\frac{\frac{\nabla P_e}{en_e} \times \mathbf{B}}{B^2}$$

If the electrons remain isotropic but have a baroclinic closure there will be a residual thawing that remains. In this circumstance the rate of flux decay due to this effect will overpower coulomb resistive rates whenever (15c) competes with the curl of the resistive emf. This condition occurs when

$$\frac{\nabla \times \frac{\nabla \cdot \mathbf{P}_e}{en_e}}{\nabla \times \eta \cdot \mathbf{J}} > 1 \qquad (16a)$$

and is satisfied in the baroclinic equation of state if

$$1 > \Xi \equiv \frac{\eta ecB}{4\pi kT_e} \qquad (16b)$$

Using the Spitzer form for $\eta = 1.15x10^{-4}z \ln \Lambda T(\text{eV})^{-\frac{3}{2}}$ (16b) becomes

$$\Xi \equiv 8.2x10^{-10}\frac{\ln \Lambda}{22}\frac{B}{50\gamma}\left(\frac{100\text{eV}}{T(\text{eV})}\right)^{\frac{5}{2}} \qquad (16c)$$

Isocontours of $\Xi(T_e, B)$ are illustrated in Figure 2. The asterisk denotes typical conditions at the noon magnetopause. We thus conclude that any baroclinic electron equation of state would thaw flux much faster than any effects associated with classical resistivity down to a misalignment from barotropic behavior of $10^{-10}radians$!

If the equation of state is isotropic but baroclinic, then whenever $\Xi < 1$ there are ambipolar dominating terms present in the plasma. For an anisotropic equation of state where $\nabla \cdot \mathbf{P}_e$ is magically parallel to the electron density then the anisotropic condition will contribute a thawing rate that overpowers resistivity if

$$\frac{P_{e,||} - P_{e,\perp}}{P_e} > \Xi$$

If there is no such barotropic like degeneracy then whenever $\Xi < 1$ is a regime where ambipolar effects overpower resistivity. The isotropic polytrope regime guarantees that the equation of state is barotropic, the ambipolar thawing contribution vanishes and resistivity controls the thawing by postulate.

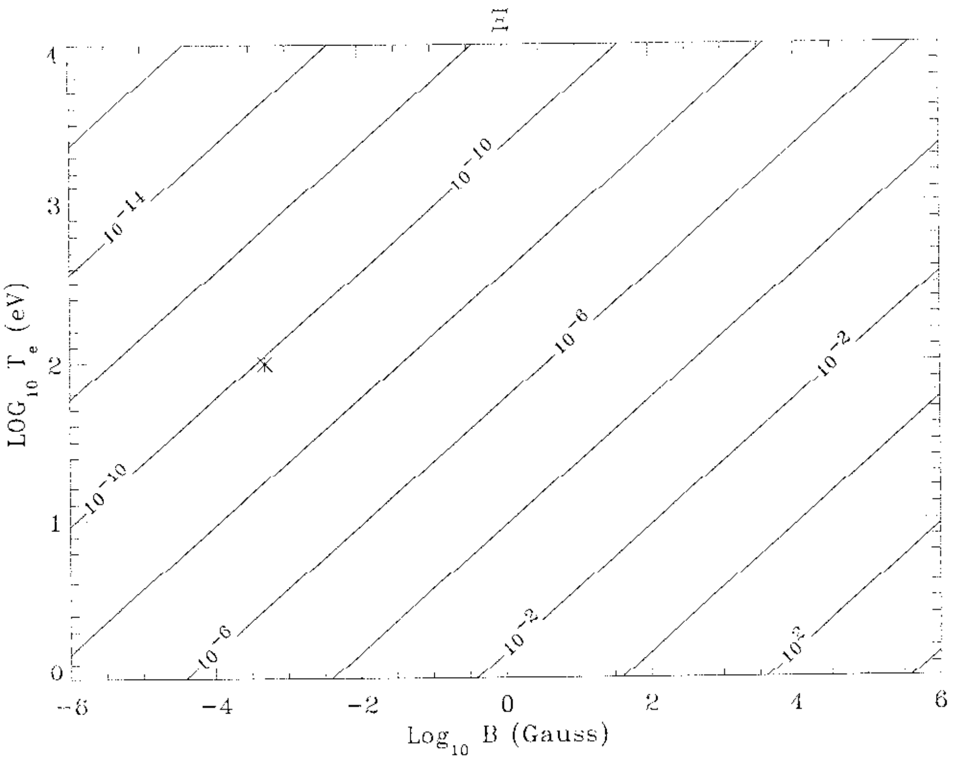

Figure 2. Isocontours of the quantity Ξ that determines the importance of ambipolar contributions to thawing of magnetic flux as a function of electron temperature in eV and magnetic field strength in Gauss. Typical nose magnetopause conditions are indicated with the asterisk.

Alternatively, if the electron pressure tensor is anisotropic, the two terms in (15b) will not in general cancel. Accordingly, in the electron bulk frame there is no flux preserving description to be foreseen in general. If a situation had a high symmetry even anisotropic plasma might conspire to make the first term of (15b) zero by geometry. The finite anisotropy of the second term spelled out explicitly in (15d) still gives the ambipolar term the edge over the resistive term at the magnetopause and at most other locations in astrophysical plasmas. An idea of the relative importance of resistive and anisotropic ambipolar effects in the thawing of magnetic flux is found by noting that the second term of (15d) scales with the pressure anisotropy so that whenever

$$\frac{|P_{\perp,e} - P_{\|,e}|}{P_e} > \Xi \tag{16d}$$

then anisotropic ambipolar thawing will overpower the more conventional resistive diffusion. Therefore, in the reconnection layer whenever the fractional pressure anisotropy exceeds 10^{-10}, resistive thawing is unimportant.

We therefore conclude that if the electrons remain magnetized at the magnetopause with a sheath magnetic field of 50γ and electron temperature of 100eV, the plasma in this region is completely at the mercy of electrons for the thawing rates either from the baroclinic-like term that does not vanish with isotropy or from the emf's associated with the very modest observed electron anisotropies ($\Xi \simeq 0.1$) of the magnetosheath. Thus it must be concluded that it is inconceivable that coulomb resistivity plays a role in the reconnection layer physics at the magnetopause or magnetotail.

There are empirical *in situ* estimates of effective resistivities in astrophysical plasmas. Estimates of the effective resistivity have been made by fitting magnetic field profiles to diffusive solutions (Burlaga and Scudder, 1974) of the Sweet-Parker type discussed below and from the energy equation and magnetic ramp scale within a collisionless shock (Scudder et al., 1986). Both methods yielded (cgs) estimates for effective resistivities of

$$\eta_{\perp,in\ situ} \simeq (0.1 - 1)\omega_{pe}^{-1} = (1 - 10)1.7 \times 10^{-6} \left(\frac{20/\mathrm{cc}}{n_e}\right)^{\frac{1}{2}} \ \mathrm{sec} \qquad (17a)$$

that should be contrasted with the numerical (cgs) Spitzer expression:

$$\eta_{Spitzer} \simeq 2.5 \times 10^{-6} \frac{\ln \Lambda}{22} \left(\frac{100\mathrm{eV}}{T_e}\right)^{\frac{3}{2}} \ \mathrm{sec} \qquad (17b)$$

If these empirical circumstance are any indication of the general scaling of anomalous resistivities, they are at worst an order of magnitude larger than the coulomb value. Even revisiting Figure 2 and increasing the threshold value of Ξ by an order of magnitude still places the observed baroclinic and or anisotropic parts of $\mathbf{P}_e$ in (15d) in complete control of magnetic flux thawing for almost all astrophysical plasmas!

Clearly, if the pressure tensor is non-gyrotropic it need not have a vanishing curl either, so the magnetic flux will not be conserved. Such effects have been foreseen to be important in the momentum balance within the diffusion region (Vasyliunas, 1975). Further details in this vein have been suggested by Dungey (1988), Lyons and Pridmore-Brown (1990, 1992). This behavior has been made more credible with a full particle code of the inner zone of a magnetic reconnection layer by Cai et al., (1994). This simulation also demonstrated the simultaneous importance of the electron and ion pressure gradient terms in (14a) within the diffusion region even though the $\frac{1}{1836}$ mass ratio would have ordinarily led the simulator to justify its neglect! In the current sheet the ions get very much hotter than the electrons and compensate for the inertial mismatch suggested by formal ordering.

In this sense observations of anisotropic or non-gyrotropic electron pressure tensors are a potential warning sign that the magnetic flux is not being preserved even in the electron's frame of reference and *the thawing of flux will not be ordered by magnetic shear or current density as it would for the case of resistive thawing.*

Even in the anisotropic case the electrons of differing energies possess different gradient and curvature drifts which when averaged over the distribution do not cancel. In this way electrons initially located on the same tube do not migrate across the magnetic field at the same rate or even in the same direction so that it is difficult (though not impossible) to guarantee that a circulation of field lines through the volume will leave the net threading flux invariant. In this sense particles once on a given tube of force find themselves at some later time on differing tubes according to their initial kinetic energy as a result of their departure from isotropy. This rearrangement makes it difficult to redefine a single field line that they were on. An analysis of this type of behavior was conducted by Scudder (1987) in the context of the fluid motion of the electron center of mass, $\mathbf{U}_e$ within a steady state normal mass flux layers, where it was shown that the only ways in the deHoffmann-Teller (1950) frame (if electrons remain magnetized) to separate the electron flow lines $\hat{\mathbf{U}}_e(x, y, z)$ from being tangent to the magnetic field lines $\mathbf{B}(x, y, z)$ while threading the layer, are electron pressure anisotropy, electron inertia and collisions.

The electron inertial terms in (14a, b) are usually small corrections and as shown below are important on smaller scales than the leading terms on the right hand sides of (14a,b). However, because they involve Maxwell equation variables in (14a) and the center of mass speed these terms are often retained without the equation of state complications mentioned above. These terms are also the only place where the mismatch of accelerations experienced by electrons and ions can play a role in determining the electric field. The recent review by Drake (1995) discusses simulations and analysis of this decoupling.

Previously Speiser (1970) and coworkers have developed the concept of inertial resistivity from these terms for use in the discussion of the neutral sheet in the magnetotail. Lyons and Speiser (1985) have suggested an effective inertial resistivity that would have the cgs order of magnitude determined by the normal component of the magnetic field, B_n, in the current sheet of the form

$$\eta_{inertial} = \frac{B_n}{n_i ec} = 3.5 \times 10^{-8} \left(\frac{B_n}{1\gamma}\right) \left(\frac{20/\mathrm{cc}}{n_i}\right) \mathrm{sec} \tag{17c}$$

which when contrasted with the coulomb resistivity implies that the inertial effect exceeds the two body resistivity effect whenever

$$\frac{\eta_{inertial}}{\eta_{coulomb}} > 1$$

which occurs when

$$\frac{B_n}{1\gamma} > 7.2 \times 10^{-9} \frac{n_i}{20/cc} \left(\frac{T_e}{100\mathrm{eV}}\right)^{-\frac{3}{2}} \frac{\ln \Lambda}{22} \tag{18}$$

that is, almost always in a space plasma!

The resistive, or $\eta \cdot \mathbf{J}$, term of (14a) is also an inertial term in disguise since η is directly proportional to the reduced mass of the electron ion pair and hence is proportional to the electron mass density as well as the electron-ion collision frequency that varies as $T_e^{-\frac{3}{2}}$. When this resistive term is theoretically included in the description of thawing frozen flux, it is invariably included as a constant η. To use such a physical or anomalous collision frequency requires knowledge of the partition of internal energy between electrons and ions to evaluate the collision frequency correctly. The anomalous resistivity models require estimates for the wave energy density as well. From the ion center of mass description of (14a), the electron temperature is unavailable except by approximating $T_e \simeq \frac{1}{2}T$ as an *ad hoc* guess. The physical resistive scale at the magnetopause is very short ($\simeq 1km$), and varies widely across any simulation domain that includes the magnetopause. Choosing η to be a constant has the practical impact of fixing the resistive scale length across the entire solution domain from the beginning of the simulation, instead of it being spatially dependent as the physical mechanism would require. If the resistive effects contribute to the variation of temperature within the sheet, the inverse dependence with T_e of the physical resistivity is not addressed by this approach. Resistive terms possess the well known therapeutic effect of numerically stabilizing MHD codes. Too often these considerations drive the simulator to retain a constant η rather than the explicit bulk parameter dependences of the resistive model considered. Often modelers explore scaling of their results with a spatially fixed resistivity until the code is no longer stable and then seek to extrapolate to the even lower resistivity regimes thought to be relevant for the geophysical problem. It should be clear that these approaches **define** the mode of interconnection to be that of *resistive thawing* without looking at other structural ways in (14) for flux to be thawed and the observed scales of current layers produced.

Two further points are emphasized by the electron form of (14b). First, the resistive emf has its origins in the collision integral between electron and ions. The $\eta\mathbf{J}$ term of (14a) is only the leading order contribution from the collision integral for the circumstance when there is bulk velocity slippage between the electron and ions considered as separate fluids. If the electrons and ion distributions are spherically symmetric in their respective frames of reference, then the form $\eta \cdot \mathbf{J}$ for the emf of (14a) is only appropriate for the *Stoke's friction regime* of the plasma. The friction between Gaussian electrons and ions is only proportional to the slippage between them and, hence the current in a collisional plasma if the relative drift speed is small and within the Stoke's regime:

$$|\mathbf{U}_i - \mathbf{U}_e| \ll \sqrt{\frac{3kT_e}{m_e}} \quad (Stokes)$$

For larger drift speeds the structure of the friction changes (Dreicer, 1959, 1960) so that the friction decreases with increasing further relative drift. This is the regime of "runaway". In a quasi-collisional plasma, collective ion acoustic instabilities can imply a change in the coefficient of resistivity and its scaling when this slippage is 42 times smaller and also depends on the $\frac{T_e}{T_i}$ ratio and wave amplitude level.

Secondly there is no guarantee, especially in low density plasmas, that the electrons and ions will remain Gaussian in their separate frames of reference while the momentum transfer is occurring, or that they are isotropic when they are not supporting a current density. All Legendre polynomials are in general required to describe *observed astrophysical velocity distribution functions in their respective* bulk velocity frames. The first order Legendre polynomial facilitates the description of heat flow about the center or mass frame of the species. The general specification of the electron-ion moment transfer contains a contribution from all odd moments of the distribution. The ohmic emf $\eta \cdot \mathbf{J}$ is the result of appropriately small mismatches in the first moments of the electrons and ions. There are additional contributions to the momentum exchange called the *thermal force emf* that **can be present whether or not there is a current density present.** This emf is *caused by the unbalanced momentum exchange that occurs between species if one* distribution is pear shaped, or *skewed*, as when a species conducts heat (Braginskii, 1965) and is usually described in the Chapman-Enskog formalism as a positive definite coefficient times the electron heat flux, $\mathbf{q}_e$. This emf produces a force on the electron that has a sense opposed to the direction of the heat flux. This function is usually ignored.

4. Significance of $\mathbf{E}_{||}$

A final form of the frozen flux theorem whenever $B > E$ is obtained by defining a special velocity field that is everywhere perpendicular to $\mathbf{B}$.

$$\mathbf{W}_E \equiv c\frac{\mathbf{E} \times \mathbf{B}}{B^2}$$

with $\mathbf{E}$ taken from the full form of the Generalized Ohm's law. The flux thawing criterion (12) takes the form

$$\frac{d\Phi}{dt} = -c\int_{S_1} \nabla \times \left(\frac{E_{||}}{B}\hat{\mathbf{b}}\right) \cdot \hat{\mathbf{n}}_1 da \tag{19}$$

Only moving with $\mathbf{W}_E$ (which does not have an evolution equation like the $\mathbf{W}$'s of (14a,b,c)) can (19) be true. (Through much of the current layer this frame is generally accessible, although when the magnetic field in the layer, B_L, is depressed relative to the ambient value in the reconnecting layer below $\frac{|U_{inflow}|}{c} \simeq 10^{-4}$ of its upstream value this frame has a superluminal speed. Thus over much of the current layer this is a useful frame to imagine keeping track of the motion of field lines.) Equation (19) can be reduced using Ampere's equation to

$$\frac{d\Phi}{dt} = -\int_{S_1} \left[c\nabla\left(\frac{E_{||}}{B}\right) \times \mathbf{B} + \frac{E_{||}}{B}\left(4\pi\mathbf{J} + \frac{\partial \mathbf{E}}{\partial t}\right)\right] \cdot \hat{\mathbf{n}} da \tag{20}$$

Consider a surface S_1 that has its normal everywhere parallel to $\mathbf{B}$, then the integral expression simplifies to

$$\frac{d}{dt}\int_{S_1(\hat{\mathbf{n}}_1=\hat{\mathbf{B}})} B da = -\int_{S_1} \frac{E_{||}}{B}\left(4\pi J_{||} + \hat{\mathbf{b}}\cdot\frac{\partial \mathbf{E}}{\partial t}\right) da \tag{21}$$

When moving with the speed $\mathbf{W}_E$ the electric field only has parallel components, so that the dot product involving the partial time derivative contains only one term. Comparing the derivation for (12) with a general $\mathbf{W}$ and the present specific $\mathbf{W}_E$ illustrates that there is no analogue of surface S_3 to be considered since by (8), $\hat{\mathbf{n}}_3 da$ vanishes. Comparing with (7) with this flow field it is now clear that (21a) may be rewritten as:

$$\int_{S_1(\hat{\mathbf{n}}_1=\hat{\mathbf{B}})} \left[\frac{1}{4\pi}\frac{\partial B}{\partial t} + \frac{1}{8\pi B}\frac{\partial E_{||}^2}{\partial t} + \frac{E_{||}J_{||}}{B}\right] da = 0 \tag{22}$$

for arbitrary enclosed area of the geometrically constrained surface S_1. Accordingly, the integrand in the $\mathbf{W}_E$ frame must vanish leaving the result that

$$\frac{\partial}{\partial t}\left(\frac{|\mathbf{B}|^2 + E_{||}^2}{8\pi}\right) = -J_{||}E_{||} \tag{23}$$

Several comments are in order about (23). First in the $\mathbf{W}_E$ frame $E_{||}$ is the entire electric field; thus the left hand side is the time rate of change of the electromagnetic energy density in this frame. For the same reason there is no Poynting flux nor a divergence in this frame. Equation (23) illustrates that the field energy density is changed only by the Joule dissipation $-\mathbf{J}\cdot\mathbf{E} = -J_{||}E_{||}$. Further if there is prexisting $E_{||}$ negative Joule dissipation need not imply erosion of magnetic flux. However, in astrophysics it is usually the case that $|\mathbf{B}| \gg |\mathbf{E}|$ which is preserved for relativistic frame transformations so that protracted Joule dissipation will usually erode $\mathbf{B}$.

The second point about (23) is the crucial role in the physics of flux erosion played by $E_{||}$. *Without $E_{||}$, conversion of magnetic field energy density cannot happen, period.* Only some of the terms of the Generalized Ohm's Law (14) can produce $E_{||}$. These are the ambipolar electric field, the resistive terms and the inertial terms and, of these, even they will not thaw flux unless they occur in the presence of parallel current flow.

5. General Causes for a Parallel Electric Field in a Plasma

The second moment of the electron kinetic equation gives a clear and complete relationship between the electric field and other properties of the plasma as indicated in (14b). The ultimate sources of $E_{||}$ are thus traceable to the imbalances of all the **parallel accelerations** experienced by the electron fluid:

$$E_{||} = -\hat{\mathbf{b}} \cdot \frac{\nabla \cdot \mathbf{P_e}}{en_e} - \frac{GM_o n_e m_e}{r^2}\hat{\mathbf{r}} \cdot \hat{\mathbf{b}}$$

$$+\frac{\hat{\mathbf{b}}}{en_e} \cdot \frac{\partial m_e n_e \mathbf{U_e}}{\partial t} + \frac{\hat{\mathbf{b}}}{en_e} \cdot \nabla \cdot (n_e m_e \mathbf{U_e U_e}) + \frac{m_e}{en_e} < \frac{\delta f}{\delta t}|_c \hat{\mathbf{b}} \cdot \mathbf{v} >_{\mathbf{U}_e} \quad (24)$$

The terms on the right hand side are the ambipolar, inertial (gravitational, accelera-tion(2 terms), and collisional) contributions, respectively to $E_{||}$. While the parallel electric fields are usually small compared to $E_\perp$, they are nonetheless exceedingly important, controlling almost all aspects of the equilibrium *along* the field. Away from the site of flux erosion the motional electric field brings in new flux that is required for a steady state reconnection pattern to form. The physics of the reconnection itself is clearly that of what causes $E_{||}$ to exist in the first place as illustrated in (23). Simultaneously, however, the reconnection is short lived if new flux is not supplied with the external dynamics outside of the inner region. While it is easiest to see what determines $E_{||}$ from (24, 14b), it should be clear where all these terms are hidden in the one-fluid form (14a): The $\eta\mathbf{J}$ comes from the collisional terms and (14a) may be recovered from (14b,c) by multiplying each of these equations by its mass density and adding, using the definition of center of mass velocity and quasi-neutrality to recover (14a).

The leading order contribution to the parallel electric field in (14a,b,or c) well away from the nearest star is almost always given by the ambipolar term:

$$E_{||} \simeq -\hat{\mathbf{b}} \cdot \frac{\nabla \cdot \mathbf{P_e}}{en_e} \quad (25)$$

which for a gyrotropic plasma reduces to the form:

$$E_{||} \simeq -\frac{1}{en_e}\left(\frac{\partial P_{||,e}}{\partial s} + \frac{P_{||} - P_\perp}{B}\frac{\partial B}{\partial s}\right) \quad (26)$$

Equation (26) shows the causal role of the parallel pressure gradients and the unbalanced mirror force along $\mathbf{B}$ in determining $E_{||}$. The ambipolar electric fields counteracts the tendency for the lighter electrons to separate from the more sluggish ions along the field; in this way the ambipolar fields are the "glue" of a quasi-neutral plasma performing a regulating role whether collisions are dominant or not.

The remaining terms in $E_{||}$ come from the inertial terms which includes the collisional terms that for a highly collisional conductor like a copper wire give the ac and dc (resistive) contributions to the traditional Ohm's law. Ordinarily, the resistive term is not important in the astrophysical context, except possibly near the surface of a star if the magnetic field is strong enough. The inertial acceleration terms are usually more important than the resistive term.

6. Relative Scale Lengths of Generalized Ohms Law Terms

Vasyliunas (1975), following Rossi and Olbert's (1970) analysis of length scales of the Generalized Ohm's law, identifies the relative scale lengths when various terms in (14 a) become important. The Hall, ambipolar, inertial and resistive correction terms to 14(a) are first important on scales lengths

$$Hall: \ ambipolar: \ inertial: \ resistive$$

$$\frac{c}{\omega_{pi}} : \beta_e^{\frac{1}{2}} \frac{c}{\omega_{pi}} : \frac{c}{\omega_{pe}} : \lambda_\eta \frac{V_A}{V}$$

and have relative sizes of

$$1 : \beta_e^{\frac{1}{2}} : \frac{1}{42} : \frac{\lambda_\eta \omega_{pi}}{c} \frac{V_A}{V}$$

where λ_η is the resistive length

$$\lambda_\eta = \frac{\eta c^2}{4\pi V_A}$$

and V is the fluid speed.

For most space plasma domains β_e is non-negligible, so that the scales for the ambipolar electric field contributions are commensurate with the ion inertial scale lengths and well separated from the electron inertial scale lengths and the resistive scales. Ambipolar contributions were crucial in the modeling of steady state, collisionless shocks. As current layers start to amplify, the ambipolar contributions to **E** with their structures different from the resistive, Hall, or inertial terms will play an important role in the subsequent evolution of the reconnection layer. Modeling reviewed by Drake (1995) and coworkers discuss the onset of kinetic Alfvén disturbances in the reconnection layer with the inclusion of such terms and that the formation of magnetic islands foreseen in resistive MHD formulation of reconnection are preempted by the inclusion of the ambipolar contributions to the generalized Ohm's law. These same authors also concluded that current layers reconnecting with inertial electron skin depth corrections were intrinsically unstable and would spontaneously break apart. Near the magnetopause as soon as currents are important with scales that make the Hall term important, so is the ambipolar term.

Winglee (1994) has started to include ambipolar effects in a large scale MHD simulations and has shown the relationships between particle code effects at the magnetopause with no counterpart in resistive MHD for example. Mandt et al., (1994) have incorporated electron pressure effects in a hybrid code modeling the reconnection layer with isotropic fluid electron pressure. The reconnection rate is found to be much larger than in resistive MHD, and is insensitive to the resistivity assumed. Hesse et al., (1995) have made the first attempts to include equations for the evolution of the electron pressure tensor in a hybrid code for reconnection. Although this is the first modeling where collisionless frozen flux violations as seen in the electron frame of reference have been inventoried, to do so has come at the expense of prohibiting the electrons from doing what they do best, conducting heat!

New effects while retaining the ambipolar terms in the baroptropic polytropic approximation have been reported by Ma and Battacharjee (1996). Abrupt increases in the current growth rate for the reconnection were noticed and a radically different spatial scaling of the non-ideal inner zone of the solution were noted. The ambipolar term in Ohm's law modified the external solution causing the resistive Sweet-Parker Y-points to collapse more toward Petschek's X configuration, but at the same time causing the reconnection rate to increase with local Alfvén speed exhausts. The parallel electric fields that evolve in such layers are much larger than the Dreicer runaway value (Ma and Bhattacharjee, private communication) so that the issues of closure comes full circle once more: the runaway contributions to the current and their contribution to the fricitional heating are incompletely inventoried in their analysis. These limitations arise from the neglect of the Stokes limitations on the friction, the constancy of the resistivity and neglect of the thermal force.

7. Levels of Magneto-Fluid-Dynamic (MFD) Approximation in the Description of Merging

Much of the early astrophysical literature of reconnection/merging involves assumptions about what simplification of the Generalized Ohm's law will be presumed adequate for the description of a physical situation. So called *Ideal MHD* approximates (14a) with (2). *Ideal-Hall MHD* entails truncating (14a) after retaining the $\mathbf{J} \times \mathbf{B}$ emf, which we have shown above is approximately like assuming field line preservation in the *electron* bulk frame. At this level there is the first possibility of separate responses of electrons and ions. Each of these *Ideal* levels of description have been further simplified and made analytically tractable in the *incompressible* limit: $\nabla \cdot \mathbf{U}_{cm} = 0$. This approximation reduces the one fluid continuity equation to

$$\frac{D}{Dt} \ln \rho = 0$$

where

$$\frac{D}{Dt} = \frac{\partial}{\partial t} + \mathbf{U}_{cm} \cdot \nabla$$

is the convective derivative. This additional approximation forestalls shock formation (without forbading their importance!). A wide and ever growing literature exists that discusses flux erosion within the framework of *Resistive (Ideal) MHD* or even *Resistive Hall MHD* which corresponds to the further retention of the $\eta \cdot \mathbf{J}$ term in the Generalized Ohms law. The electron pressure term of the Generalized

Ohm's is important in plasmas of very low density where it is usually referred to as the ambipolar electric field. In the absence of collisions the inertial terms of (14a-c) are present as well. As a group the inertial and ambipolar terms are sometimes referred to as the *"collisionless"* terms in the Generalized Ohm's law. In the magnetotail Speiser (1970) and coworkers have emphasized the role of the inertial terms in producing an effective resistivity as the particles oscillate about the current sheet.

Given the scaling arguments of the preceding section and the observations, nothing short of *anisotropic, ambipolar, Hall MHD*, more properly called Magneto-Fluid-Dynamics, makes any sense to this author for modeling of reconnection for near earth plasmas.

8. Geometrical Requirements for Reconnection

Favorable geometry is required between juxtaposed field lines so that reconnection is possible, (Hill, 1975; Cowley, 1976). Whether attempting component or anti-parallel merging, some components of B of the opposite sides perpendicular to the"X", or merging line must reverse across the current sheet so there is a current to exert a deflection force onto the plasma along reconnected field lines. The components of **B** parallel to the merging line appear to be unconstrained.

A predetermined boundary is identified to be the site of non-idealness and the plasma is pushed toward that boundary with its embedded magnetic flux by a perpendicular electric field that is tangential to the interface boundary where the system is non-deal. If the external boundary conditions are in balance with the flux dissipation and the expulsion of plasma along the current boundary, a steady state may result.

9. An Overview of Fluid Models

Analytic models of reconnection differ in the ways that the "external" inflow is arranged to interact with the inner zone. Most of the early work treats the region within this central layer as a "black box", or boundary layer, where resistive diffusion was assumed to dominate the evolution of the flux, while the external problem is arranged to (a) match a steady state with the boundary condition imposed by the large scale boundary conditions to match the (b) interface condition supposed for the inner region. Work of this type by Sweet (1958) and Parker (1957) dealt with spontaneous reconnection caused by resistive departures from Ideal Incompressible MHD: the inner diffusion region was long and narrow line of zero area of the "Y" form: $\succ\!\!-\!\!\prec$ and the inflow Alfvén mach number was very low of order $M_A = R_M^{-\frac{1}{2}}$, where the magnetic Reynold's number $R_M = \frac{LV_{a,in}}{4\pi\eta}$ is customarily very large $> 10^6$. Petschek (1964) introduced compressibility and standing slow shocks into the external solution outside of the diffusion region, with the effect of focusing the inflow region towards a narrower interval, thereby enhancing the reconnection rate to $M_A \simeq \ln R_m^{-1}$. Unlike the Sweet-Parker scenario where all the energy conversion must take place along the diffusion boundary layer, Petchek's solution achieved a more efficient conversion of magnetic to bulk energy by using the shocks to irreversably heat up and accelerate the plasma along the entire "X" interface radiating from the corners of the finite length reconnection line of the diffusion region. In this way, the reduced length of the diffusion region of his solution did not simultaneously "choke" the overall rate at which new flux could be brought up to the diffusion region along the symmetry axis. Topologically the Petschek geometry more nearly approximates an "X" than the back to back abutted "Y" geometry of Sweet-Parker. Additionally, the Petschek solution with slow shocks required a structured outflow. In Sweet-Parker and Petschek's model the inflow region was characterizes by very small Alfvén Mach numbers, while the outflow or exhaust region was a local Alfvén mach number 1 regime with the fluid flow nearly parallel to the highly kinked field lines in the exhaust region. The release of these stresses propagates at the local Alfvén speed perpendicular to the kink in the field accelerating the attached fluid in this external exhaust region to a flow velocity of this order. This accelerating outflow came to be known as "jetting" and for many synonymous with the reconnection process. In the early 1980's the lack of detection of accelerated jetting was taken as strong evidence that reconnection was not involved in mass interchange at the magnetopause.

Still within the resistive ideal MHD description, Levy, Petschek and Siscoe (1964) presented a model that addressed the asymmetric nature of the possible reconnection at the magnetopause. In this model the strong field side is treated as nearly a vacuum compared to the high density of the magnetosheath with its weaker magnetic field. Each pair of symmetrical slow shocks in the Petschek solution are now replaced with a rotational discontinuity and a slow expansion fan. The rotational discontinuity reconciles the arbitrary angles between the asymptotic sheath field and that of the magnetosphere (Petschek and Thorne, 1967) generalizing the antiparallel case of Levy et al. and Sweet-Parker. The tangential stress repartition made by the RD is the subject of several observational tests of reconection discussed below.

Priest and Forbes (1986) succeeded in identifying a family of resistive ideal MHD reconnection solutions which contained those of Petschek(1964) and Sonnerup (1970) as special cases of the inflow field to the reconection layer: Sonnerup's (1970) solution is the first place in the one parameter variation of B across the diffusion region that permitted smooth steady state flow with all flow lines deflected from the inflow symmetry line, whereas Petschek's solution is the last solution where all stream lines converge towards the inflow symmetry axis. Unlike Sweet-Parker all flow lines in this family of solutions do **NOT** go through the diffusion region. A significant fraction of the inflowing plasma is deflected and accelerated at the standing waves/ shocks in the external solution. In addition new solutions were also found including the limiting case for this family beyond which a steady solution can no longer be found.

The inclusion of a normal component at a previously defined tangential discontinuity has been used by Heyn et al., (1988) and Biernat et al., (1989) to develop a more detailed, if approximate picture of the reconnection layer, as a solution to the so-called Riemann initial value problem. These authors outline the set of nested slow shocks, expansion or rarefaction fans that are predicted to stand inside of two flanking Alfvén or rotational discontinuities. The Alfvén wave on the magnetosheath side is suggested to carry most of the current at that reconnection layer, but propagates at a lower speed along the magnetopause than does the Alfvén discontinuity along the magnetoshere side of the reconnection layer. Detailed patterns for the changes in the fluid parameters through such a layer were computed and the effects of non-zero tangential flow speeds at the magnetopause away from the nose were included. Recent work by Lockwood et al., (1996) have explored the consequences of this structured layer withing the reconnection layer for the time of flight dispersed ions in the boundary layers of the cusp.

10. Overview of Single Particle Descriptions

Many of the results of the reconnection theory follow from enforcing conservation laws. Single particle discussions without using the fluid equations achieve similar predictions essentially by invoking the same conservation laws that transcend details of thermal physics description whether it is fluid or collisionless. This type of work has established that the resistivity caused by binary collisions is not the essential ingredients for reconnection. The fluid equation description provides a reduced level of description for examining the consequences of the conservation equations.

Hill (1975) modeled totally collisionless behavior and obtained reconnnection rates similar to those of the MHD theory using the kinetic equations and their conservation laws. He found that particles streaming into the field reversal region must carry enough momentum to exactly balance the stress in the magnetic field in steady state. Alternatively "...the particle motion in the field reversal must provide exactly the electric current required by the field geometry through Ampere's equation." The self consistency requirement is that the jump of the total tangential stress vanish, viz:

$$\Delta[\int (m_e f_e + m_i f_i) v_x v_k d^2 v - B_x B_k/(4\pi)] = 0$$

where $\hat{\mathbf{x}}$ is the normal to the field reversal plane, $\mathbf{k}$ is the direction of the component of $\mathbf{B}$ that reverses across the current sheet. This can be easily rephrased into the Walén condition for a rotational discontinuity, or for a shock as another example of a possible imbedded normal mass flux layer. Quite generally it is concluded that if the incoming pressure is isotropic, then the merging speed is much less than V_A and is the local Alfvén speed times the small angle between the incoming field lines and the field reversal plane, consistent (within factors of 2) with that of fluid theory. From the particle picture merging ceases at the fire hose limit. For a given field the merging speed is independent of the magnitude of the initial plasma pressure *so long as it is isotropic*. The merging speed is reduced if $P_{||} > P_{\perp}$ since this situation reduces the magnetic tension required to accelerate and eject the plasma from the reversal region. For given asymptotic field strengths B_1 and B_2 the merging electric field has a maximum value for antiparallel fields and decreases monotonically to zero as the angular separation of the fields decreases. Flow patterns similar to Petschek's were achieved. The limitations of this model are: (i) the analysis is only valid so long as the particle drift in the current sheet layer perpendicular to the "plane of the X" is small compared to the characteristic dimensions of the system (Hill's analysis was performed sliding along the current in deHoffmann-Teller frame which is supposed to exist for all the orbits in question); and (ii) was only a fully collisionless treatment; particles only interact with the large scale electromagnetic field and not with each other; dissipation of electro- magnetic energy is dominated by inertial effects rather than collisional Joule heating.

Cowley (1980, 1982, 1995), Lockwood (1995) and coworkers (Lockwood et al., 1996) have discussed the form for the phase space distributions after they have interacted with the Alfven wave and rotational discontinuity within the reconnecting layer assuming that the magnetic moment of reflected and transmitted particles is conserved. This type of analysis has recently been reviewed by Fuselier (1995) and Onsager (1997). To date these models have not attempted to solve for the accessability of the populations from the opposite sides of the layer, nor have they addressed the problem of maintaining quasi-neutrality in the reconnection layer.

11. Signatures and Tests of Magnetic Reconnection

There are several robust tests for a normal mass flux discontinuity in MHD. They are, however, usually predicated on a locally planar magnetopause that is time stationary in its own frame of reference which may not be realized in practice.

Finding a normal component B_n of **B** at the magnetopause is an important clue about the site and occurrence of reconnection. The operational problem of finding small components perpendicular to a *free surface* that is oriented in an unknown direction and moving with an unknown velocity relative to the moving spacecraft has left us with a paultry few instances when the normal component is defensibly non-zero and steady through the layer. This situation is required for a planar normal mass flux rotational discontinuity. Numerous techniques have been developed with the best examples found in Sonnerup and Ledley (1979).

If the magnetopause has a normal mass flux, then the plasma streaming toward the magnetopause has in its external regime well removed from the magnetopause current layer a unipolar motional electric field $\mathbf{E_T}$ parallel to the magnetopause. From equation (2) this electric field (even when unmeasured) can be estimated using plasma and magnetic field data; this approach is especially accurate when the plasma is sampled near the magnetopause but outside of the precursors associated with the magnetopause current layer. Alternatively, long wire probes can be used to determine this vector, but in the sub-solar magnetopause, $\mathbf{E_T}$ is primarily in the GSE z direction and, until the recent Polar mission, has not been directly sampled, but synthesized from spin plane antennae under the assumption that $\mathbf{E}\cdot\mathbf{B} = 0$. The free surface problem alluded to in the case of the magnetic field normal component is also an issue here. With the electric field measurements there is a new complication relative to the Galilean invariant magnetic field vectors: Estimates of $\mathbf{E}$ relevant to the observer on the discontinuity are not the same electric field measurements recorded on the the spacecraft since: (1) the magnetopause is moving and (2) the spacecraft is also moving. Accordingly the tranformation of electric fields between frames will require knowledge of the magnetic field and the relative vector velocity of the free surface as viewed from the spacecraft. Accordingly, to use these measurements care must be taken to model and remove the effects of the magnetopause motion and acceleration. New techniques have recently been devised in this connection and have been reviewed by Sonnerup (1995). Often $\mathbf{E}$ is not directly measured, but inferred using (2) and even then without the knowledge of the center of mass velocity. It is always prudent to keep the observables and the inferences clearly identified in this detective work!

A variant of determining $\mathbf{E_T}$ are tests that contrast observational consequences of both $\mathbf{E_T} \neq 0$ and $B_n \neq 0$ *without determining the local magnetopause normal or relative velocity.* The first approach exploits the well known Walén (1944) jump condition for a rotational discontinuity

$$\Delta \mathbf{U}_{cm} = \alpha \Delta \mathbf{B}, \tag{27a}$$

which relates the **changes** in the asymptotic center of mass velocities of the plasma to the changes in the magnetic field over the same interval. As a vector relation this condition is actually three separate conditions with a common, specified proportionality constant (cf Hudson, 1970)

$$\alpha = \frac{1}{4\pi\rho}(1 - \frac{P_{\|}}{P_{\perp}}). \tag{27b}$$

This relationship results immediately from tangential stress balance. Note carefully that the condition as stated requires a determination of the plasma center of mass velocity and obtains only for the asymptotic states well removed from currents of the RD that implement the angular rearrangement of the asymptotic magnetic fields.

Until recently most of the "smoking gun" evidence for reconnection from *in situ* observations has come with the performance of numerous tests to verify the Walén test for a rotational discontinuity (cf. Paschmann et al., 1979, and Sonnerup et al., 1981). The data used for these tests came from ISEE 1 and 2 ion measurements of energy per unit charge analyzers that did not differentiate ions by their charge to mass. Moments of the observed fluxes at their observed energies were performed assuming all fluxes were protons (Paschmann private communication, 1996). Follow on studies (e.g. Sonnerup et al., 1990) with Ampte-IRM also used energy per unit charge analysis, achieving better time resolution by performing moments on board the spacecraft prior to transmission. Many layers were found with good correlation between the right and left hand sides of the Walén relationship. The agreement was however not perfect, with examples being dubbed "good" when the Walén expectation based on the magnetic field changes were corroborated at the greater than the 50% level. There were also some very good examples with nearly perfect correspondence. Attempts have been made (Paschmann, 1997) to organize the incidence of "good" Walén intervals with size of the magnetic shear to no avail, their being excellent Walén matches at low or high shear with no clear preference. The one clear ordering appears to be with plasma magnetosheath β. Whenever $\beta < 5$ Walén matches appear to be uniformly good or better.

In order to improve the leverage on the Walén comparisons, data through the current carrying layers have been used (Sonnerup et al., 1995) even though the jump conditions as stated by Walén are not supposed to work there. This follows from the fact that the Walén condition can be derived from the conservation of $\mathbf{E}_T$ across the rotational discontinuity. From 14a, there are Hall modifications to the emf that imply that the ideal MHD result for the asymptotic regimes is not strictly applicable within the layers. Although $\mathbf{E}_T$ is conserved if the layer is one dimensional and time stationary, the components of (14a) that are required to evaluate $\mathbf{E}_T$ include terms other than those ion terms of (2) that have been used in this data evaluation. These data should be resorted based on current density, not just shear angle to see if the cases that failed the Walén test were those of higher than nominal current density or, equivalently, thinner than those that fit the expectation well. Alternatively, the Walén condition is fullfilled (at the MHD sizes of measurable E) throughout the current carrying layer provided electron bulk velocity measurements are used to estimate E for this study (Scudder, 1984). Perhaps these will be available from the upcoming Cluster instrumentation. The role of alpha particles in the determination of the center of mass velocity from an E/Z analyzer warrants further consideration. Finally, in view of the discussion above of the Generalized Ohm's law, it may be that current density is not the index for thawing of magnetic flux at work at the magnetopause. Seeking such an ordering tacitly assumes that the resistive term of (14a) is the controlling factor at the magnetopause, even though the current carrying layers are much thicker than such an organization would imply. Another candidate for the possible organization of the good and not good Walén tests as indexes of magnetic "thawing" would be the size of the magnetosheath electron pressure anisotropy.

As compelling as the Walén tests are for detecting potential reconnection layers away from the diffusion region, there are circumstances where the Walén condition can be fulfilled without there being a rotational discontinuity there. Except for forbading a normal component, the flow speeds are unconstrained across a tangential discontinuity. If the flow fields were arranged to meet the Walén condition in the transverse components and the normal component changes were unchanged but small, this is a permitted class of tangential discontinuities that if present would signal a locally closed magnetopause. This diagnostic degeneracy is more than an academic one: precisely this type of "mistake" was made when counting the frequency of detection of rotational discontinuities in the free flowing solar wind (cf. Burlaga, 1971 and Belcher and Davis, 1971). In fact the early counts for rotational discontinuities decreased markedly when this was realized. Accordingly, arguing that chance matches of the Walén condition are improbable had to be reworked! From a theoretical point of view tangential discontinuities are the non propagating limit of the slow branch, so that preexisting flows with tangential stress balance might be organized previously by the evolution of slow wave fronts into local tangential discontinuities leaving organized flows on either side of local tangential discontinuities. It is also known that high beta is not favorable to slow mode wave propagation; as slow expansion waves are required in the asymmetrical reconnection scenario, the observed beta organization of the "good" Walén data at weak shear may have their origin in the flows on either side of the layer that have temporarily closed down to become a tangential discontinuity.

It should also be stated for consideration that the mere certification of a rotational discontinuity in a given locale does *not* provide incontrovertible evidence of ongoing reconnnection. The rotational discontinuity is one of a number of MHD discontinuities found in magnetohydrodynamics that are required in general to support arbitrary boundary conditions. As an example rotational discontinuites and slow shocks are observed in the free flowing solar wind (Burlaga, 1995) without any reflexive implication of their implying ongoing reconnection. These structures are swept over and through the magnetopause region and will be modified at the bow shock. Nevertheless, there is by this conunterexample no one to one association with if its a rotational discontinuity satisfying Walenś conditions, then there must be ongoing reconnection. The rotational discontinuity is for almost all geometries necessary, but not sufficient to indicate frozen flux violation.

Important for removing chance confirmations of reconnection are multiple, ancillary pieces of information about the magnetic topology, as from the leaking of particles from the magnetosphere along a separatrix (Sonnerup et al., 1981; Scudder et al., 1984), or fluid speed in the magnetopause frame equal to the local Alfvén speed, tests of conservation of energy. Many of these test have been performed as in the paper by Sonnerup et al., (1981) and the review by Paschmann (1984) for example, and often they also point to a picture of witnessing the reconnection layer away from the separator line implicit in the asymmetric Levy-Petschek-Siscoe (1964) picture.

In the noon equatorial magnetopause geometry of opposed fields brought together normal to the current sheet, the expectations of "jetting" within the reconnection layer, or more precisely the lack of same, was once used as an argument against ongoing reconnection at the magnetopause. The initial surveys for reconnection flows from ISEE were indexed by enhanced flow speeds up against the magnetopause, consistent with Petschek's picture in the antiparallel geometry. There are, however, other tests for normal mass flux through the magnetopause performed by Aggson and co-workers that were not biased by looking for speed enhancements. Using measurements of the DC electric and magnetic field data, Aggson et al.,(1983, 1984, 1985) asked whether there was a Galilean frame shift to the deHoffmann-Teller frame that could null the detectable MHD sized motional electric field components even though $\mathbf{E}$ was varying across the current carrying layer of the magnetopause. Not only was this technique successful, this type of analysis rather routinely suggested that the magnetopause was open, since the existence of the deHoffmann-Teller frame, except in bizarre counterexamples, cf. Paschmann , (1985), is indicative of a normal component of $\mathbf{B}$ and of $\mathbf{E}_T$ to the free surface since $|\mathbf{V}_{HT}| = c\frac{|\mathrm{ET}|}{B_n}$. Aggson et al. studies (1984) further illustrated examples of open rotational shear layers where the fluid speed decreased rather than the field aligned "jetting" increase of the Levy et al., (1964) paradigm. Scudder (1984) looked at the general problem of a rotational discontinuity standing in the sheath flow and mapped out the general circumstances where net speed increases like Petschek were expected and where the speed decreasing solutions were expected. When the external tangential sheath flow is comparable or exceeds the Alfvén speed at the local density, there is an increased incidence with rotational shear that the reconnected flux will have a reduced speed in the earth's frame even while the flow vector acceleration still meet the Walén condition.

Care should be taken to note that outside the current layers the electron bulk speed and that of the ion plasma would be equal. The Walén tests are then the plasma counterpart of conservation of tangential $\mathbf{E}_T$. Scudder (1984) used (14b) from the electrons to see what kind of behavior should be expected for electrons (as a field line tracer) through the layer. Scudder also illustrated the expected variation of the electron fluid speed within the rotational shear layer, explicitly using (14b) as a more appropriate picture of the statement of frozen flux through the current carrying layer. Finally he delineated the frequency of occurrence of aligned flow within the reconnection layer. It is clear that strictly anti-parallel fields yields specialized field aligned jetting that are not the most general expectation. Accordingly, the many observations of speed going down across a boundary layer at the magnetopause are **NOT** first line evidence that reconnection has not locally populated those layers. Also, jetting need not be the exhaust flow geometry, especially away from the subsolar point.

12. Summary

The theoretical description of magnetic reconnection in the plasmas of geospace is conceptually and computationally challenging. This regime is not like the fusion regime and it is not a regime controlled by two body resistivity or even anomalous resistivities. The most important needed improvement beyond always retaining the Hall emf's in the description of reconnnection is the inclusion of the dominant terms in the Generalized Ohm's Law that control the parallel electric field. These terms are by all accounts the anisotropic ambipolar emf and possibly the inertial finite Larmor radius (FLR) ion gyroradius effects that mock up resistivity in the weakest part of the current sheets. Curiously, all of these effects are coherent in character as opposed to the stochastic resistive terms that are now in common use for this type of modeling. This situation has some hauntingly parallel aspects to the evolution of our understanding of collisionless shocks over the past 20 years. There a decided paradigm shift has occurred away from a stochastic two body or anomalous resistivity dominated scenario to a more realistic, but complicated picture with numerous coherent effects facilitating flux dispersal and thermalization. For reconnection, the challenging non-gyrotropic ambipolar terms are usually not included because they require a frontal assault on the issues of pressure equations of state and closure issues at the energy equation level that are amongst the most challenging problems in astrophysics. Nevertheless, these terms in the Generalized Ohm's law, as estimated by observations, control the erosion of magnetic flux in almost all realms of astrophysics. Finally, a truly predictive theory of reconnection for geospace plasmas must involve some approach to assess the evolution of the full electron pressure tensor, so that its curl is properly assessed, rather than ignored in the dynamical estimates of the erosion of magnetic flux. Such an approach must also retain the physics of electron heat conduction that was omitted from the promising start in this direction made by Hesse et al., (1995). In the intervening time the observations may further strengthen this impression by documenting the variation of electron pressure anisotropy and departures from gyrotropy in and near reconnection layers and possibly within the diffusion region. It is possible that these overlooked signatures of thawing of magnetic flux will better organize layers that are otherwise thought to be candidates for being open such as by "passing" Walén or related tests.

References

Aggson, T.L., Gambardella, P.J. and Maynard, N.C.: 1983, 'Electric field measurements at rotational magnetopause discontinuities', *J. Geophys. Res.* **88**, 10,000.

Aggson, T.L., Gambardella, P.J., Maynard, N.C. Ogilvie, K.W.,and Scudder, J.D.: 1984, 'Observations of plasma deceleration at a rotational discontinuity', *Geophys. Res. Lett.* **11**, 8.

Aggson, T.L., Gambardella, P.J. and Maynard, N.C.: 1985, 'Reply', *J. Geophys. Res.* **90**, 7636.

Alfvén, H.: 1942, 'On the existence of electromagnetic-hydrodynamic waves', *Ark. Mat.* **29B**, 2.

Axford, W. I.: 1984, in E.W. Hones, Jr., (ed.). ' Magnetic field reconnection', *Magnetic Reconnection in Space and Laboratory Plasmas*, Geophys. Monongraph 30, AGU, Washington, 1ff.

Biernat, H.K., Heyn, M.F., Rijnbeek, R.P., Semenov, V.S., and Farrugia, C.J.: 1989, 'The structure of the reconnection layer: application to the earth's magnetopause', *J. Geophys. Res.* **94**, 287.

Braginskii, S.I.: 1965, 'Transport process in a plasma', in *Reviews of Plasma Physics*, 1, Consultant's Bureau, New York.

Belcher, J.W. and Davis, L. Jr.: 1971, 'Large amplitude Alfvén waves in the interplanetary medium 2.', *J. Geophys. Res.* **76**, 16.

Burlaga, L.F.: 1971, 'Nature and origin of directional discontinuities in the Solar Wind', *J. Geophys. Res.* **76**, 4360.

Burlaga L.F. and Scudder, J.D.: 1974, 'Sweet's mechanism in the solar wind', *Ap. J.* **191**, L149.

Burlaga L.F.: 1995, 'Interplanetary Magnetohydrodynamics', Oxford University Press, New York.

Cai, H.J., Ding, D.Q., and Lee, L.C.: 1994, 'Momentum transport near a magnetic X line in collisionless reconnection', *J. Geophys. Res.* **99**, 35.

Cowley. S.W.H.: 1976. 'Comments on the merging of nonparallel magnetic field'. *J. Geophys. Res.* **81**, 3455.

Cowley, S.W.H.: 1980, 'Plasma populations in a simple open magnetosphere', *Space, Sci. Rev.* **26**, 217.

Cowley, S.W.H.: 1982, 'The causes of convection in the Earth's magnetosphere: a review of developments during the IMS', *Rev. Geophys. and Space Phys.* **20**, 3, 531.

Cowley, S.W.H.: 1995, in P. Song, B.U.Ö. Sonnerup, and M.F. Thomsen, (eds.), 'Theoretical perspectives at the magnetopause', *Physics of the Magnetopause*, Geophys. Mono 90, AGU Washington, p 29.,1995.

deHoffmann, F., and Teller, E.: 1950, 'Magnetohydrodynamic shocks', *Phys. Rev.* **80**, 692.

Dreicer, H.: 1959, 'Electron and ion runaway in a fully ionized gas, I', *Phys. Rev.* **115**, 238.

Dreicer, H.: 1960, 'Electron and ion runaway in a fully ionized gas, II', *Phys. Rev.* **117**, 329.

Drake, J.F.: 1995, P. Song, B.U.Ö. Sonnerup, and M.F. Thomsen, (eds.), 'Magnetic reconnection: a kinetic treatment', *Physics of the Magnetopause*, Geophys. Mono 90, AGU Washington, 155.

Dungey, J.W.: 1988, 'Noise free current sheet', in *Proc. of an International Workshop in Space Plasma*, ESA SP-285, 15.

Fainberg, J., Osherovich, V.A., Stone, R.G., MacDowall, R.J., and Balogh, A.:1996, in D. Winterhalter, J.T. Gosling, S. R. Habbal, W.S Kurth, (eds.) ,'Ulysses observations of electron and proton components in a magnetic cloud and related wave activity. *Solar Wind Eight.* AIP press. Woodbury, NY, p 554.

Fuselier, S.A.: 1995, P. Song, B.U.Ö. Sonnerup, and M.F. Thomsen, (eds.), 'Kinetic aspects of reconnection', *Physics of the Magnetopause*, Geophys. Mono 90, AGU Washington, 181.

Hesse, M., Winske, D. and Kuznetsova, M.: 1995, 'Hybrid simulations of collisionless reconnection in two-dimensional current sheets', *J. Geophys. Res.* **100**, 21,815.

Heyn, M.F., Biernat, H.K., Rijnbeek, R.P., and Semenov, V.S.: 1988, 'The structure of the reconnection layer', *J. Plasma Phys.* **40**, 235.

Hill, T.W.; 1975, 'Magnetic merging in a collisionless plasma', *J. Geophys. Res.* **80**, 4689.

Hudson, P.D.:1970, 'Discontinuities in an anisotropic plasma and their identification in the solar wind', *Planet. Space, Sci.* **18**, 1611.

Kleva, R.G., Drake, J.F., and Waelbroecke, F.L.: 1995, 'Fast reconnection in high temperature plasmas', *Phys. Plasmas* **2**, 23.

Lee, L.C.: 1995, P. Song, B.U.Ö. Sonnerup, and M.F. Thomsen, (eds.), 'A review of magnetic reconnection: MHD Models', *Physics of the Magnetopause*, Geophys. Mono 90, AGU Washington, p 139.

Levy, R.H., Petschek, H.E., and Siscoe, G.L.: 1964, 'Aerodynamic aspects of the magnetospheric flow', *AIAA J.* **2**, 2065.

Lockwood, M.: 1995, 'Location and characteristics of the reconnection X-line deduced from low-altitude satellite and ground-based observations: 1 theory', *J. Geophys. Res.* **100**, 21,791.

Lockwood, M., Cowley, S.W.H. and Onsager, T.G.: 1996, 'Ion acceleration at both the interior and exterior Alfvén waves associated with the magnetopause reconnection site: signatures in cusp precipitation', *J. Geophys. Res.* **101**, 21,501.

Lottermoser, R. -F, and Scholer, M.: 1997, 'Undriven magnetic reconnection in magnetohydrodynamics and Hall magnetohydrodynamics', *J. Geophys. Res.* **102**, 4875.
Lyons, L.R. and Speiser, T.W.: 1985, 'Ohm's law for a current sheet', *J. Geophys. Res.* **90**, 8543.
Lyons, L.R. and Pridmore-Brown, D.C.: 1990, 'Force balance near an X line in a collisionless plasma', *J. Geophys. Res.* **95**, 20,903.
Lyons, L.R. and Pridmore-Brown, D.C.: 1992, 'Force balance near an X line with $\mathrm{E} \cdot \mathrm{J} < 0$', *J. Geophys. Res.* **97**, 2955.
Ma, Z., and Bhattacharjee, A.: 1996, 'Fast impulsive reconnection and current sheet intensification due to electron pressure gradients in semi-collisional plasmas', *Geophys. Res. Lett.* **23**, 1673.
Mandt, M.E., Denton, R.E., and Drake, J.F.: 1994, 'Transition to whistler mediated reconnection', *Geophys. Res. Lett.* **21**, 73.
Mozer, F.S., Torbert, R., Fahleson, U.V., Falthammar, C.G., Gonfalone, A., and Pedersen, A.: 1978, 'Electric field measurements in the solar wind, shock, magnetosheath, magnetopause and magnetosphere', *Space Sci. Rev.* **22**, 791.
Onsager, T. G. and Lockwood, M.: 1997, 'High latitude particle precipitation and its relationship to magnetospheric source regions', *Space Sci. Rev.*, this volume.
Osherovich, V.A., Farrugia, C.J., Burlaga, L.F., Lepping, R.P., Fainberg, J., and Stone, R.G.: 1993, 'Polytropic relations in interplanetary clouds', *J. Geophys. Res.* **98**, 15331.
Osherovich, V.A., Farrugia, C.J., and Burlaga,L.F.: 1995, 'The non-linear evolution of magnetic flux ropes: 2 Finite beta plasma', *J. Geophys. Res.* **100**, 12307.
Parker, E.N.: 1957, 'Sweet's mechanism for merging magnetic fields in conducting fluids', *J. Geophys. Res.* **62**, 509.
Paschmann, G., Sonnerup, B.U.Ö., Papamastorakis, I., Sckopke, N., Haerendel, G., Bame, S.J., Asbridge, J.R., Gosling, J.T., Russell, C.T., and Elphic, R. C.: 1979, 'Plasma acceleration at the earth's magnetopause: evidence for reconnnection', *Nature* **282**, 243.
Paschmann, G.: 1984, in E.W. Hones, Jr., (ed.), 'Plasma and particle observations: implications for reconnection', *Magnetic Reconection in Space and Laboratory Plasmas*, Geophys. Monongraph 30, AGU, Washington, 114.
Paschmann, G.: 1985, 'Comment on "Electric field measurments at the magnetopause I: observations of large convective velocities at rotational magnetopause discontinuities" by T.L. Aggson, P.J. Gambardella and N.C. Maynard', *J. Geophys. Res.* **90**, 7629.
Paschmann, G.: 1997, 'Observational evidence for transfer of plasma across the magnetopause', *Space Sci. Rev.*, this volume.
Petschek, H.E.: 1964, in W.N. Hess, (ed.),'Magnetic field annihilation', *The Physics of Solar Flares*, NASA SP-50, Washington DC.
Petschek, H.E. and Thorne, R.M.: 1967, 'The existence of intermediate waves in neutral sheets', *Ap. J.* **147**, 1157.
Priest, E.R. and Forbes, T.G.: 1986, 'New models for fast steady state magnetic reconnection', *J. Geophys. Res.* **91**, 5579.
Rossi, B. and Olbert, S., *Introduction to the Physics of Space*, McGraw Hill Book Co., New York, 1970.
Scudder, J.D. and Olbert, S.: 1979, 'A theory of local and global processes which affect solar wind electrons, 2. experimental support', *J. Geophys. Res.* **84**, 6603.
Scudder, J.D. Ogilvie, K.W., and Russell, C.T.: 1984, in E. Hones, (ed.), 'The relationship of flux transfer events to magnetic reconnection', *Magnetic Reconnection in Space and Laboratory Plasmas*, AGU Geomonograph Series 30, Washington, DC, 153-154.
Scudder, J.D.: 1984, 'Fluid signatures of rotational discontinuities at the Earth's magnetopause', *J. Geophys. Res.* **89**, 7431-7440.
Scudder, J.D., Mangeney, A., Lacombe, C., Harvey, C., and Aggson, T.L.: 1986, 'The resolved layer of a collisionless, high β, super-critical, quasi-perpendicular shock wave: 2: dissipative fluid electrodynamics', *J. Geophys. Res.* **91**, 11053.
Scudder, J.D.: 1987, 'The field aligned flow approximation for electrons within layers possessing a normal mass flux: a corollary to the deHoffmann - Teller theorem', *J. Geophys. Res.* **92**, 13447-13455.

Scudder, J.D.: 1992, 'On the cause of temperature change in inhomogeneous, low-density astrophysical plasmas', *Ap. J.* **398**, 299.

Siscoe, G.L.: 1983, in R.L. Carovillano and J.M. Forbes (eds.), 'Solar system magnetohydrodynamics', *Solar Terrestrial Physics*, p 11-100, D. Reidel., Amsterdam.

Sittler, E.C., Jr. and Scudder, J.D.: 1980, 'An empirical polytrope law for solar wind thermal electrons between 0.45 and 4.76 AU: Voyager 2 and Mariner 10', *J. Geophys. Res.* **85**, 5131.

Sonnerup, B.U.Ö.: 1970, 'Magnetic re-connexion in a highly conducting incompressible fluid', *J. Plasma Phys.* **4**, 161.

Sonnerup, B.U.Ö. and Ledley, B.G.: 1979, 'OGO 5 magnetopuase structure and reconnection', *J. Geophys. Res.* **84**, 399.

Sonnerup, B.U.Ö., Paschmann, G. , Papamastorakis, I., Sckopke, N., Haerendel, G., Bame, S.J., Asbridge, J.R., Gosling, J.T., and Russell, C.T.: 1981, 'Evidence for magnetic field reconnection at the earth's magnetopause', *J. Geophys. Res.* **86**, 10,049.

Sonnerup, B.U.Ö.: 1984, in E.W. Hones, Jr., (ed.),' Magnetic reconnection at the magnetopause: a review', *Magnetic Reconection in Space and Laboratory Plasmas*, Geophys. Monongraph 30, AGU, Washington, 92.

Sonnerup, B.U.Ö., Papamastorakis, I., Paschmann, G., and Luhr, H.: 1990, 'The magnetopause for large magnetic shear: analysis of convection electric fields',*J. Geophys. Res.* **95**, 10,541.

Sonnerup, B.U.Ö., Paschmann, G., Phan, T. D.: 1995, in P. Song, B.U.Ö. Sonnerup, and M.F. Thomsen, (eds.), 'Fluid aspects of the reconnection layer', *Physics of the Magnetopause*, Geophys. 90, AGU Washington, 167.

Speiser, T.W.: 1970, 'Conductivity without collisions or noise', *Planet. Space Science* **18**, 613.

Sweet, P. A.: 1958, in B. Lehnert, (ed.), 'The neutral points of solar flares', *Electromagnetic Phenomena in Cosmical Research*, Cambridge University Press, London.

Vasyliunas, V.M.: 1972, 'Nonunniqueness of magnetic field line motion', *J. Geophys. Res.* **77**, 6271.

Vasyliunas, V.M.: 1975, 'Theoretical models of magnetic field line merging: 1', *Rev. Geophys. and Space Physics.* **13**, 303.

Walén, C.: 1944, 'On the theory of sunspots', *Ark. Mat. Astron. Fys.*, 30A.

Winglee, R.M.: 1994, 'Non-MHD influences on the magnetospheric current system', *J. Geophys. Res.* **99**, 13437.

Address for correspondence: J. D. Scudder, Department of Physics and Astronomy, University of Iowa, Iowa City, IA 52242, USA