VII. MAGNETIC RECONNECTION and CORONAL EVOLUTION

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MAGNETIC RECONNECTION ON THE SUN

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ABSTRACT. Magnetic reconnection is a fundamental process in astrophysics which plays many different roles on the Sun. The classical theory of fast steady reconnection due to Sweet, Parker, Petschek and Sonnerup has been unified in a theory which possesses many new regimes that depend on the boundary conditions at large distances. For example, the *flux pile-up regime* possesses diverging flows, a long central diffusion region and a reconnection rate that is much larger than the Petschek value. Recent numerical experiments however, often possess three features that are not present in the earlier theoretical models, namely highly curved inflow field lines, separatrix jets and reversed current spikes, and so an attempt is described to include these features in a new theoretical model.

Several solar phenomena where reconnection is believed to be operating are described. Cancelling magnetic features in photospheric magnetograms are probably evidence of reconnection submergence. Most prominences are of inverse polarity and have upflows driven by reconnection below them. Coronal heating may well be due to reconnection in many small current sheets. In solar flares reconnection may be driven below an eruptiing active-region prominence and hence power the high-temperature flare loops and chromospheric ribbons.

1. Introduction

It was realised 30 years ago that the energy for a flare can only come from the coronal magnetic field and that it is being released very fast - at a significant fraction of the Alfvén speed - and so this was the stimulus for the discovery of fast reconnection mechanisms by Sweet, Parker and Petschek. The emphasis was on steady mechanisms because the overall energy release continues for many thousands of Alfvén times, although of course it is modulated in a time-dependent manner.

In a vacuum, reconnection is a trivial process, but, in a plasma atmosphere such as the Suns, normally the plasma is attrached very effectively to the magnetic field. It is only where the magnetic gradients are, say, a million times stronger than normal that the magnetic field can slip through the plasma and reconnect (Figure 1). There are three important effects of such a process. Firstly, the global topology of the magnetic field may be changed since the connectivity of the field lines may be altered. For example, initially in Figure 1 the point A is joined to point B and finally it is connected to point C. This may affect the behaviour of particles and heat which tend to travel along the field lines. Secondly, inflowing stored magnetic energy is converted into heat, bulk kinetic energy and fast particle energy. Thirdly, reconnection creates large electric currents, electric fields, shock waves, filamentation, each of which may be involved in the acceleration of fast particles.

E. R. Priest and V. Krishan (eds.), Basic Plasma Processes on the Sun, 271–291.

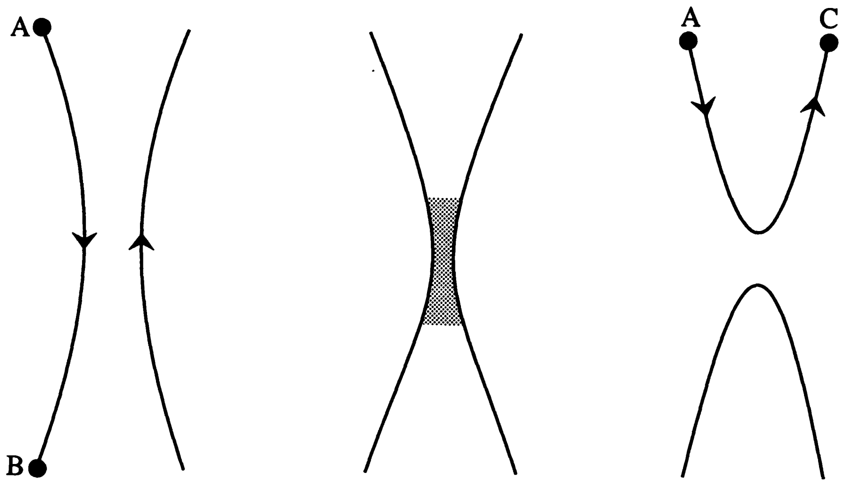

Figure 1. The breaking and reconnecting of field lines in a region (shaded) of very strong magnetic gradient.

Now reconnection is realised to be important in a wide range of solar phenomena, such as coronal heating, coronal mass ejections, magnetoconvection, dynamo generation, prominences, and bright points, as well as the many types of solar flare. We have seen in the talks by Heyvaerts, Hollweg and Van Ballegooijen how reconnection in many current sheets may provide the heating of the solar corona. Thus all

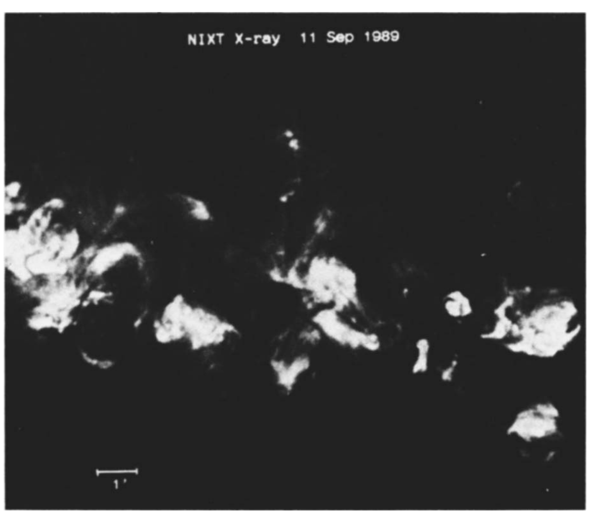

Figure 2. X-ray picture of the solar corona (courtesy L Golub)

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of the beautiful 3/4 arcsec structure in the spectacular X-ray picture from Golub's Normal Incidence X-ray Telescope on the Sept 11 (1989) rocket flight may be a direct consequence of reconnection. The X-ray bright points that are often present in coronal X-ray pictures, although strangely not in Figure 2, are now thought to be produced by magnetic reconnection, though not always by emerging flux. Reconnection also takes place below most prominences and releases the energy in a solar flare, creating both the Hα ribbons and the hot flare loops after an active-region prominence has erupted. Ahead of an erupting prominence one sees a coronal mass ejection (Figure 3), a huge magnetic bubble which is sometimes observed to disconnect from the Sun.

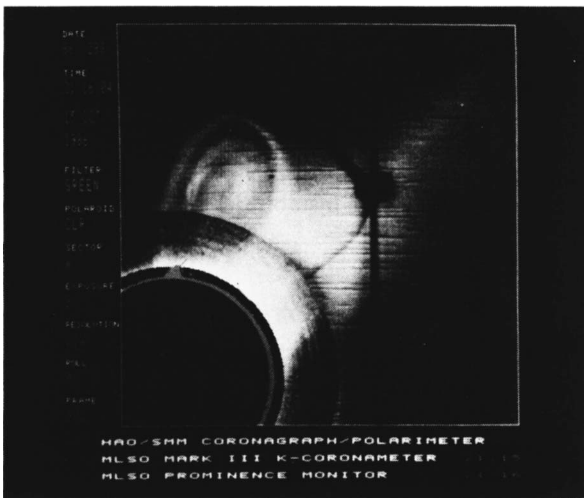

Figure 3. A coronal mass ejection (courtesy A Hundhausen)

In magnetohydrodynamics the plasma velocity (**v**) and magnetic field (**B**) are primary variables, with the electric current

$$\mathbf{j} = \nabla \times \mathbf{B}/\mu \approx B/(\mu L) \tag{1.1}$$

and electric field

$$\mathbf{E} = - \mathbf{v} \times \mathbf{B} + \mathbf{j}/\sigma \tag{1.2}$$

deducible from them if required. (This is quite different from laboratory electromagnetism, where **E** and **j** are primary.) The flow and field are determined by the equation of motion

$$\rho \frac{d\mathbf{v}}{dt} = - \nabla p + (\nabla \times \mathbf{B}) \times \mathbf{B}/\mu, \tag{1.3}$$

and the induction equation

$$\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{v} \times \mathbf{B}) + \eta \nabla^2 \mathbf{B} \tag{1.4}$$

where $\eta = (\mu\sigma)^{-1}$ is the magnetic diffusivity. Equation (1.4) implies that the magnetic field changes in time due to two effects. The first term on the right represents the advection of magnetic field lines with the plasma. The second represents diffusion through the plasma and at the same time magnetic energy is converted to heat ohmically at a rate j^2/σ, where j is given by (1.1).

The time-scale for magnetic dissipation is given by equating the first and third terms in (1.4), namely

$$\tau_d = \frac{L^2}{\eta} \approx 10^{-9} L^2 \; T^{3/2}, \tag{1.5}$$

which is enormously long (10^{14} seconds) for a typical global coronal length-scale ($L = 10^7$m) and coronal temperature ($T = 10^6$K). Thus, in order to release magnetic energy, one needs to create extremely small length-scales in sheets or filaments and therefore very large magnetic gradients and electric currents (Eqn 1.1). This may be done in three ways: as propagating sheets in shock waves; as stationary sheets in magnetic equilibria near X-points or separatrices; and as a result of instabilities such as the tearing mode or the coalescence instability.

2. Theory of Fast Magnetic Reconnection

2.1 CLASSICAL THEORY

A region of oppositely directed or sheared magnetic field may be linearly unstable to the breaking and reconnection of field lines by the tearing mode instability. When the tearing mode develops nonlinearly or when a current sheet forms dynamically or when magnetic sources are driven together, then a quasi-steady state of fast nonlinear reconnection may be reached. The classical models are sketched in Figure 4. The Sweet-Parker (1958) model is a simple diffusion region. In the Petschek mechanism (1964) the diffusion region occupies only a small central location, while most of the energy conversion occurs at standing slow-mode shock waves that accelerate and heat the plasma to form two hot fast outflowing streams. In Sonnerup's (1970) model an extra set of discontinuities is standing in the flow ahead of the slow shocks, but Vasyliunas (1975) pointed out that these are slow-mode expansion waves which need to be generated externally at discrete points in the flow which would not be present in astrophysical applications.

Sonnerup and Priest (1975) discovered an exact solution of the nonlinear MHD equations in which an incompressible stagnation point flow

$$v_x = -\frac{V_0 x}{a}, \qquad v_y = \frac{V_0 y}{a}$$

satisfying $\nabla \cdot \mathbf{v} = 0$ carries in oppositely directed (but straight) magnetic field lines. The equation of motion (1.3) for steady flow determines the pressure

$$p = \text{const} - \frac{1}{2}\rho v^2 - B^2/(2\mu)$$

and for a field $B(x)\hat{y}$ Ohm's Law becomes

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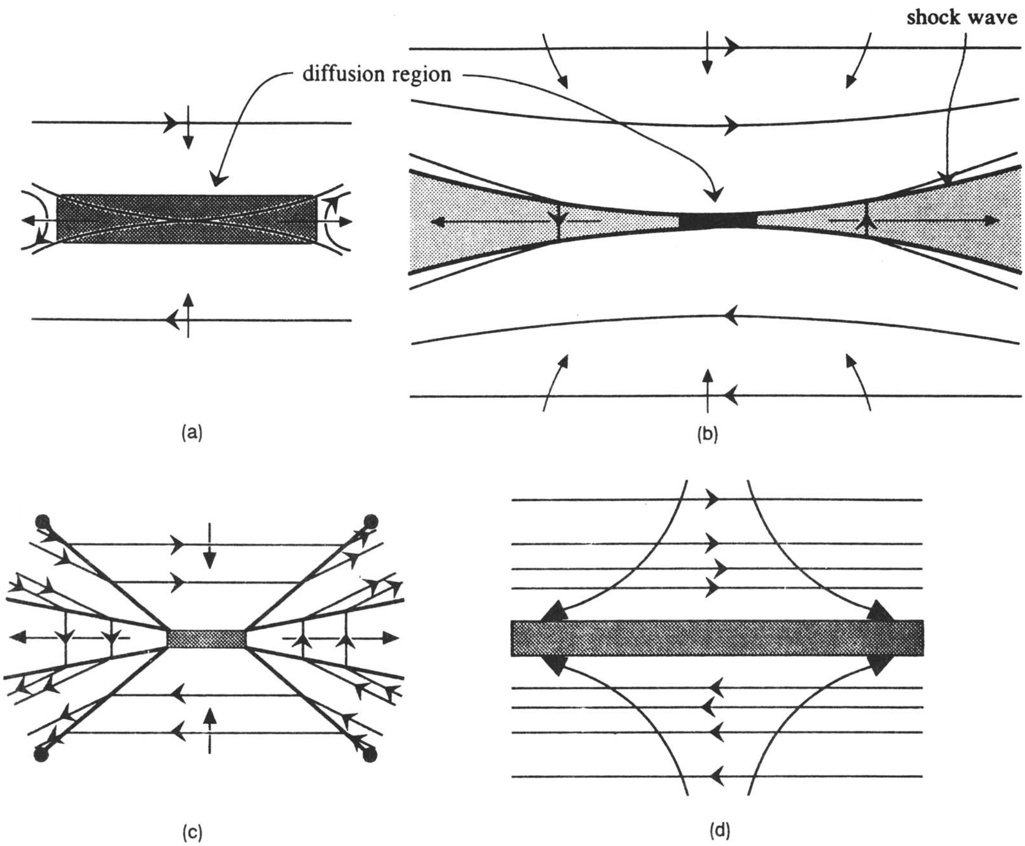

Figure 4. The classical energy conversion mechanisms of (a) Sweet-Parker, (b) Petschek, (c) Sonnerup and (d) Sonnerup and Priest.

$$E_0 - \frac{V_0 x}{a} B = \eta \frac{dB}{dx} \tag{2.1}$$

where the electric field (E_0) and flow scale (V_0) are constant. The solution of (2.1) at large distances has the field (B) increasing like x^{-1} as it is carried in, but eventually at a distance of order $(\eta a/v_0)^{1/2}$ the diffusion term on the right becomes important and causes B to decrease to zero at the origin. This solution has been generalised to three-dimensions and also, by Gratton *et al*. (1988) to include viscous flows.

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In the Sweet-Parker diffusion region of length 2L and width 2ℓ one simply writes down order of magnitude relations between the inflow speed (v_i) and the outflow speed (v_{Ai}), which is the Alfvén speed based on the inflow magnetic field (B_i). Thus for a steady state the plasma flows in at the speed

$$v_i = \frac{\eta}{\ell} \tag{2.2}$$

at which the magnetic field is trying to diffuse outwards. Also the conservation of mass into and out of the region gives

$$Lv_i = \ell v_{Ai}, \tag{2.3}$$

Thus eliminating ℓ between (2.3) and (2.4) gives in dimensionless form the inflow Alfvén Mach number (i.e. the reconnection rate) as

$$M_i = \frac{1}{R_{mi}^{1/2}} \tag{2.4}$$

where $M = v/v_A$ and $R_m = Lv_A/\eta$ is the magnetic Reynolds number. In practice R_m is typically 10^6-10^{12} and so the reconnection is very slow ($M_i \sim 10^{-3}$-10^{-6}), which was why Petschek sought a faster mechanism to explain energy release in a flare.

Petschek's analysis was disarmingly simple. The magnetic field decreases substantially from a uniform value (B_e) at large distances to a value B_i at the entrance to the diffusion region, while the flow speed increases from v_e to v_i. The object is to determine for a given B_e the maximum value of v_e (in dimensionless form $M_e = v_e/v_{Ae}$). The effect of the shocks is to provide a normal field component B_N which essentially produces the distortion in the inflow field from the uniform value B_e at large distances. Thus, if the inflow field is potential, the distortion may be regarded as being produced by a series of monopole sources along the x-axis between $|x| = L$ and $|x| = L_e$, say. The result is that, as the diffusion region is approached the field strength falls to

$$B_i = B_e - \frac{2}{\pi} \int_L^{L_e} \frac{2B_N}{x} dx$$

or

$$B_i = B_e - \frac{4B_N}{\pi} \log \frac{L_e}{L}. \tag{2.5}$$

At the shock waves the condition that they be standing is

$$\frac{B_N}{\sqrt{(\mu\rho)}} = v_e,$$

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and Petschek estimates the maximum reconnection rate (M_e^*) by putting $B_i = \frac{1}{2} B_e$ in (2.5) to be

$$M_e^* = \frac{\pi}{8 \log R_{me}} .$$

In practice this is typically 0.01, much greater than the Sweet-Parker rate.

2.2 UNIFIED THEORY OF FAST STEADY ALMOST-UNIFORM RECONNECTION

Vasyliunas (1975) clarified the physics of Petschek's mechanism by pointing out that the inflow region has the character of a fast-mode expansion with the pressure and field strength decreasing and the flow converging as the magnetic field is carried in. A fast-mode disturbance has the plasma and magnetic pressure increasing or decreasing together, while a slow-mode disturbance has the plasma pressure changing in the opposite sense to the magnetic pressure. An expansion makes the pressure decrease while a compression makes it increases, even in the incompressible limit. Sonnerup's model possesses slow-mode expansions that are unlikely because they are discrete. Vasyliunas suggested that a Sonnerup-like solution may be possible with a diffuse slow-mode expansion spread throughout the inflow region, making the field strength increase, the pressure decrease and the flow diverge as the field lines are carried in, although he was unable to find such a solution.

I wanted to understand Vasyliunas's distinction mathematically and was also puzzled at many strange features of some of the numerical reconnection experiments such as much longer diffusion regions than Petschek, diverging flows and large pressure gradients. Also what is the relation to the stagnation point flow solution? Can a Sonnerup-like solution be found without the extra discontinuities? And can a model in a finite region be produced, like the numerical experiments?

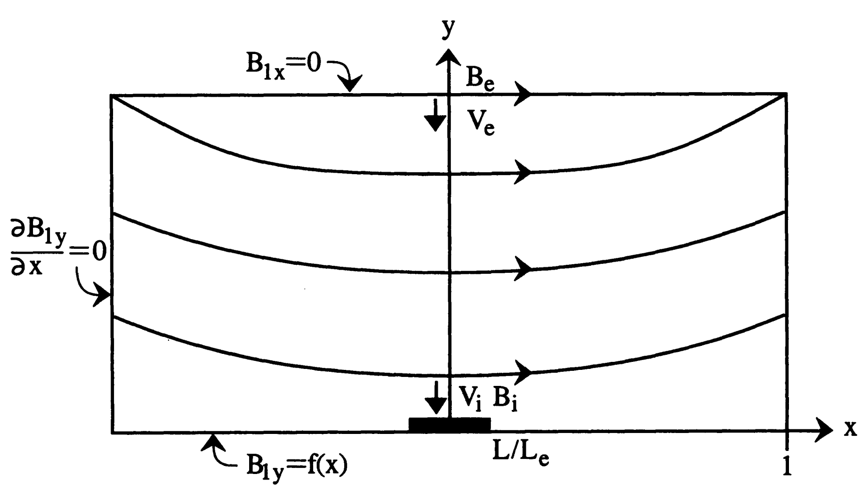

Figure 5. Notation and boundary conditions.

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During one summer Terry Forbes and I (Priest and Forbes, 1986) tried to answer such questions by seeking fast, steady almost-uniform reconnection solutions to the equations for two- dimensional, incompressible flow, namely

$$\rho(\mathbf{v}\cdot\nabla)\mathbf{v} = -\nabla p + (\nabla\times\mathbf{B})\times\mathbf{B}/\mu \tag{2.6}$$

$$\mathbf{E} + \mathbf{v}\times\mathbf{B} = \mathbf{0} \tag{2.7}$$

$$\nabla\cdot\mathbf{v} = \nabla\cdot\mathbf{B} = 0. \tag{2.8}$$

The solutions are almost-uniform in the same sense as Petschek's, namely that one performs a linear expansion about a uniform field

$$\mathbf{B} = B_0\hat{\mathbf{x}} + \mathbf{B}_1 + \ldots$$

$$\mathbf{v} = \mathbf{v}_1 + \ldots .$$

Neglecting the pressure gradient (2.6) then reduces to a potential field

$$\nabla^2 A_1 = 0 \tag{2.9}$$

where

$$B_{1x} = \frac{\partial A_1}{\partial y}, \qquad B_{1y} = -\frac{\partial A_1}{\partial x} .$$

The solution to (2.9) subject to the boundary conditions that

$$B_{1x} = 0 \quad \text{on} \quad y = L$$

$$\frac{\partial B_{1y}}{\partial x} = 0 \quad \text{on} \quad x = L$$

$$B_{1y} = 0 \quad \text{on} \quad x = 0$$

$$B_{1y} = f(x) = \begin{cases} 2B_N\, x/L & 0 \le x \le L \\ 2B_N & L \le x \le L_e \end{cases}$$

is

$$A_1 = -\sum_0^{\infty} \frac{a_n}{\left(n+\frac{1}{2}\right)\pi} \cos\left[\left(n+\frac{1}{2}\right)\pi\frac{x}{L}\right] \cosh\left[\left(n+\frac{1}{2}\right)\pi\left(1-\frac{y}{L}\right)\right] \tag{2.10}$$

where

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$$a_n = \frac{4B_N \sin\left[\left(n + \frac{1}{2}\right)\pi L/L_e\right]}{L/L_e\left(n + \frac{1}{2}\right)^2\pi^2 \cosh\left[\left(n + \frac{1}{2}\right)\pi\right]}.$$

This represents a Petschek-type solution with a weak fast-mode expansion. From (2.7) the first-order flow ($\mathbf{v}_1 = (E/B_0)\hat{\mathbf{y}}$ is uniform but the second-order flow is converging. By calculating B_i one can deduce a relation between M_e and M_i, which shows that, as Petschek had expected, M_e does indeed possess a maximum value which is close to Petschek's estimate.

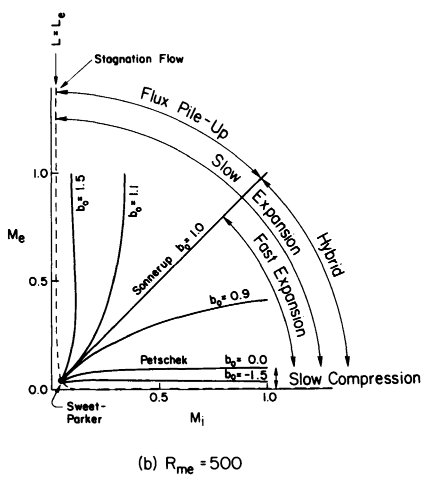

Figure 6. Inflow Alfvén Mach number M_e as a function of M_i for several values of b.

As well as putting Petschek's analysis on a firmer basis, one can, however, generalise the theory by including a first-order pressure gradient, which simply changes (2.9) to

$$\nabla^2 A_1 = -\frac{\mu}{B_0}\frac{dp_1}{dy}$$

and the solution (2.10) to

$$A_1 = -\sum_0^\infty \frac{a_n}{\left(n + \frac{1}{2}\right)\pi}\left(by - \cos\left[\left(n + \frac{1}{2}\right)\pi\,\frac{x}{L}\right]\cosh\left[\left(n + \frac{1}{2}\right)\pi\left(1 - \frac{y}{L}\right)\right]\right)$$

When b = 1, the inflow field on the y-axis is uniform and we have found the Sonnerup-like solution with a slow-mode expansion across the whole inflow region. However, there are many other solutions for the other values of the parameter b, which is determined by the nature of the flow on the inflow boundary, since the horizontal flow speed at the corner $(x,y) = (L_e,L_e)$ is proportional to $(b - 2/\pi)$. As b increases so the inflow turns from being

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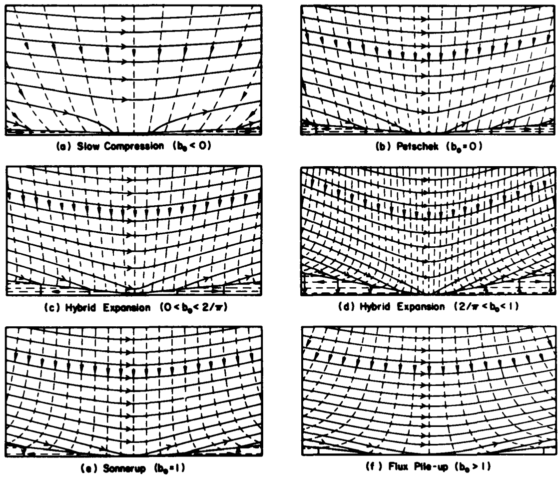

Figure 7. Magnetic field lines and streamlines for several members of the unified theory.

converging (and therefore producing slow-mode compressions) to being diverging (with strong slow-mode expansions). The latter comprise a flux pile-up regime with long diffusion regions. The way that the reconnection rate (M_e) varies with M_i and b is shown in Figure 6, where it is seen to be faster than the Petschek rate for regimes with $b > 0$.

The main results from the above analysis are that the type of reconnection regime and the rate of reconnection depend sensitively on the inflow boundary conditions, with the Petschek (b=0) and Sonnerup-like (b=1) solutions being particular members of a much wider class. Jardine and Priest (1988) have recently extended the theory to include higher orders, compressibility and energetics. Also, it has been compared with a variety of numerical experiments (Forbes and Priest, 1987).

2.3 NONUNIFORM RECONNECTION WITH SEPARATRIX JETS

Numerical experiments such as that shown in Figure 7 reveal four puzzling features that are not present in the classical models of reconnection:

(i) different types of inflow;
(ii) separatrix jets;
(iii) reversed current spikes;
(iv) highly curved field lines in the inflow region.

The first feature is the one which Priest and Forbes (1986) attempted to model with their unified almost-uniform theory. The other three features have been tackled in a new nonuniform theory by Priest and Lee (1990).

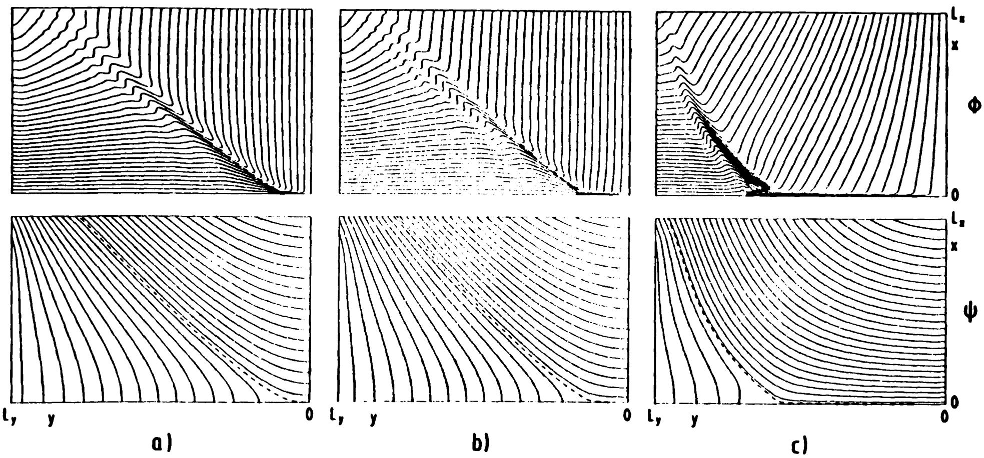

Figure 7. Numerical experiment on reconnection (Biskamp, 1986)

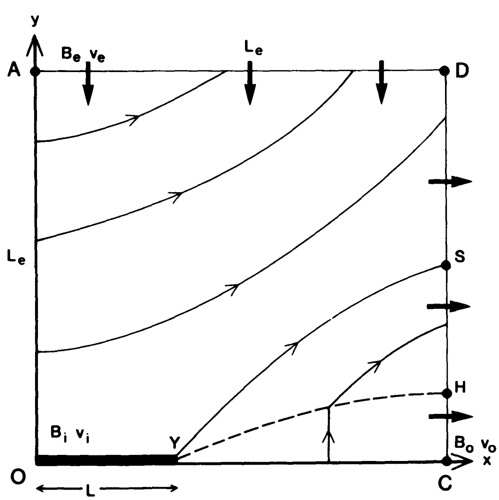

Figure 8. The notation for reconnection with a highly curved inflow.

Feature (iii), the reversed current spikes at the ends of the diffusion region, slow down the streams of plasma that are emerging from the diffusion region and partially divert them along the separatrix jets. This is a consequence of imposing boundary conditions at the outflow boundary that give a mismatch with the outflow from the diffusion region.

By contrast, feature (iv), the highly curved inflow field lines and the associated wide shock angle, are a direct result of the form of the inflow boundary conditions. In general the number of such conditions that can be imposed equals the number of MHD characteristics that are propagating information into the region. For our case of two-dimensional, sub-Alfvénic, incompressible, essentially ideal flow, there are three imposed conditions. For instance, if one prescribes

$$v_x = 0, \quad v_y = \text{constant}, \quad p = \text{constant} \tag{2.11}$$

on the inflow boundary (AD), then the MHD equations (2.6)-(2.8) imply that

$$B_x = \text{constant}, \quad B_y = 0,$$

so that the straight field lines are carried in by a uniform flow without curving and reconnection is impossible. If, on the other hand, boundary conditions only slightly different from (2.11) are imposed, then reconnection with weakly curved inflow field lines may be produced. Thus it is entirely reasonable to expect that conditions greatly different from (2.11) could produce a highly curved inflow.

Now equations (2.8) may be satisfied identically by writing **v** and **B** in terms of a stream function (Ψ) and flux function (A), namely

$$v_x = \frac{\partial \Psi}{\partial y}, \quad v_y = -\frac{\partial \Psi}{\partial x}, \quad B_x = \frac{\partial A}{\partial y}, \quad B_y = -\frac{\partial A}{\partial x}.$$

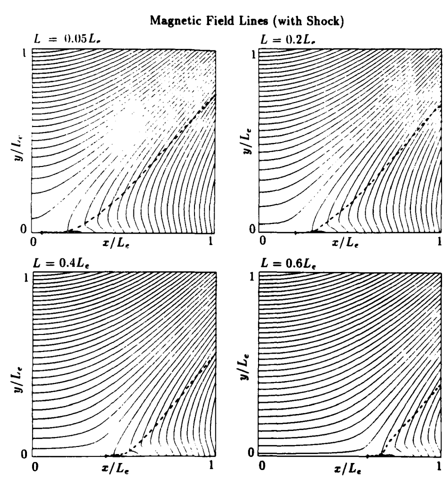

Figure 8. Field lines for one quadrant of nonuniform reconnection.

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Then the problem of nonuniform reconnection with an imposed inflow (v_e) and field strength (B_e) on the inflow boundary producing a diffusion region (OY) of half-length L with a separatrix YS and slow shock YH may be tackled in three steps.

First of all, in the upstream region ahead of YH suppose for simplicity that both the plasma speed (v) and sound speed are much smaller than the Alfvén speed (v_A). Then (2.6) implies that j = 0 and so for a potential field with a current sheet we may use complex variable theory to pick

$$B_y + iB_x = B_i\left(\frac{z^2}{L^2} - 1\right)^{1/2} \tag{2.12}$$

where z = x+iy and there is a cut (a current sheet) from z = -L to z = L. Then (2.7) implies that the flow velocity may be deduced from

$$\Psi = v_e B_e \int \frac{ds}{B} \tag{2.13}$$

where the integral is along a field line.

The second step is to calculate the position of the shock from the characteristic curve

$$\Psi + A = \text{constant}$$

that passes through the end point (Y) of the diffusion region. Then the shock relations are applied to deduce the conditions just downstream of YH. Finally, one needs to solve the MHD equations in the downstream region subject to the appropriate boundary conditions at the shock and at the outflow boundary CH. In general these equations may be written

$$\mathbf{v}\cdot\nabla A = -v_e B_e \tag{2.14}$$

$$\mathbf{v}\cdot\nabla \omega = \mathbf{B}\cdot\nabla j \tag{2.15}$$

where $\omega = -\nabla^2\Psi$ is the vorticity.

For example, the results of assuming $v >> v_A$ and so taking $\omega = 0$ or

$$\nabla^2\Psi = 0$$

with Ψ imposed along the boundary YHC and Ψ = constant on YC are shown in Figures 8 and 9. The shock (strictly speaking an Alfvénic discontinuity of slow-mode compressional type in this incompressible model) is shown dashed but is rather weak and has little effect on the magnetic field. The effect of the reversed current spike downstream of the diffusion region shows up in the field lines of abnormal curvature and in the spreading of the streamlines. Also the separatrix jet is prominent and makes streamlines follow the separatrix as they pass through it. As the current sheet decreases in length, so the inflow speed increases up to a value that depends on the inflow Alfvén Mach number. Results have also been obtained by solving the full equations (2.14) and (2.15).

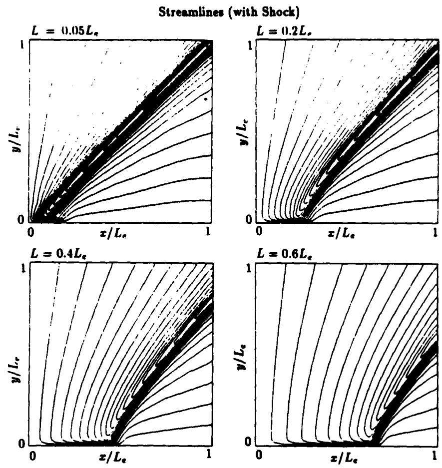

Figure 9. Streamlines corresponding to Figure 8.

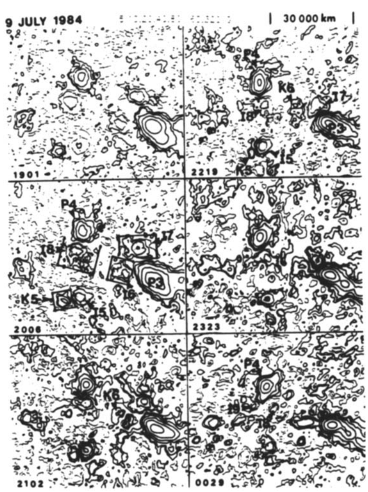

Figure 10. Cancelling magnetic features shown by rectangles in a set of photospheric videomagnetograms (Martin et al, 1985)

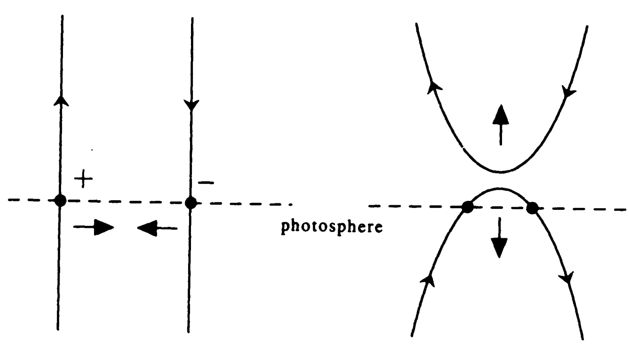

Figure 11. Reconnection submergence as an explanation for the cancellation of magnetic fragments.

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3. The Role of Reconnection on the Sun

3.1 CORONAL HEATING IN CURRENT SHEETS

There may well be several mechanisms at work in heating different parts of the corona. Almost certainly X-ray bright points, for example, are caused by reconnection, though not always, as was believed until recently, by emerging flux. They certainly occur above pairs of oppositely directed magnetic fragments, but Martin (1986) has demonstrated with a videomagnetograph that such fragments are often cancelling rather than emerging (Figure 10). It is my view that, as such fragments come together, their magnetic field lines reconnect in the atmosphere above, creating one loop which moves out while another submerges - such a process may be called *reconnection submergence* (Priest, 1987). Such an explanation has natural theoretical grounds and would in a natural manner create the X-ray bright points, subflares, macrospicules and minifilaments that are often observed to be associated with cancelling magnetic features.

Reconnection could also be responsible for energy release in current sheets created if smooth magnetic equilibria do not always exist when photospheric footpoint motions are complex (Parker, 1972). Mikic *et al* (1988) have recently performed a 3D ideal MHD numerical experiment to show how such current sheets could be created. They assume the pressure is zero, the flows are much slower than the Alfvén speed, and they adopt a $64 \times 64 \times 64$ point mesh. The footpoints of an initially straight uniform magnetic field are

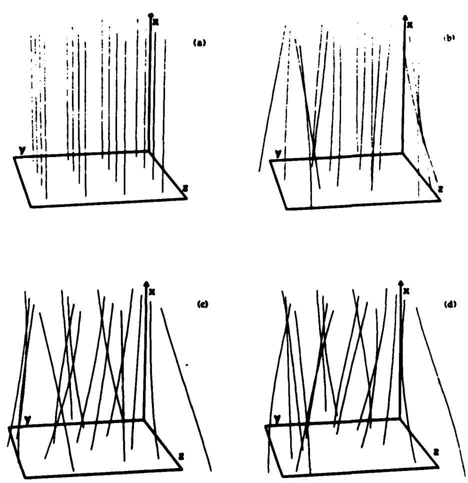

Figure 12. Field lines at intervals of $40\tau_A$ (Mikic *et al*, 1988)

braided with random footpoint motions. A series of smooth equilibria are produced but there is a transfer to small scales with filamentary electric currents that grow like $\exp(0.01t/\tau_A)$.

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3.2 PROMINENCES AND GLOBAL FLUX BALANCE

Quiescent prominences may well lie at the boundaries of giant cells and most have Inverse Polarity, so that the magnetic field threads through the prominence in the opposite direction from what one would expect on the basis of the observed photospheric line of sight components. The observed upflows of plasma in prominences on the disc have therefore

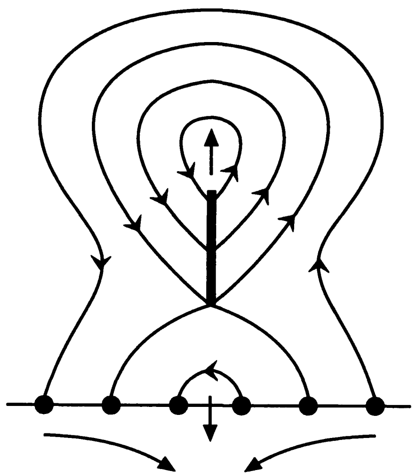

Figure 13. A prominence of Inverse Polarity with flux cancellation

been suggested to be created by magnetic reconnection below the prominence in response to converging motions of the photospheric footpoints (Malherbe and Priest, 1983; Priest, 1987). In addition such a process of flux cancellation below a prominence would be consistent with Sara Martin's observations and would build up helical structure in the prominence magnetic field (Van Ballegooijen and Martens, 1989) and so create the structure that forms the basis of the new Twisted Flux Tube Model of prominences (Priest *et al*, 1989)

An important question is how the global magnetic flux balance in the corona is maintained. Magnetic buoyancy brings flux to the surface over a wide range of scales, but then what happens to it? On a large scale flux is lost in coronal mass ejections at a rate of about 0.5 per day carrying a flux of 10^{21}Mx. When the field in a coronal mass ejection becomes stretched out, a disconnection is needed eventually to prevent an indefinite build-up of flux in the interplanetary medium although it is rarely observed, perhaps because it usually occurs after the mass ejection fades from view. On a small scale we have seen that cancelling magnetic features often create X-ray bright points and perhaps small plasmoids (and HRTS jets) are associated with such events. The plumes that overlie bright points in coronal holes (Ahmad and Webb, 1978; Holt and Mullan, 1986) may have an outflow of 100 km s^{-1}, which would be enough to supply the high-speed solar wind. Furthermore, a flux escape of 10^{21}Mx per day would provide the heat required for the active-region corona, while 10^{20}Mx per day would be enough for a quiet region.

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3.3 RECONNECTION IN SOLAR FLARES

Reconnection plays several roles in solar flares. Small flares may be created by the interaction of neighbouring flux systems, either when new flux emerges or when existing bipolar flux moves laterally or when separate magnetic elements cancel by reconnection submergence. Large flares involve three phases: the slow preflare rise of a prominence; the flare onset and impulsive phase when reconnection begins at many points below the prominence; the main phase when quasi-steady reconnection continues and creates hot coronal loops and bright chromospheric ribbons as the opened-out field closes back down. The role of reconnection in a large (two-ribbon) flare is therefore to release the energy in the impulsive and main phases and, in some cases, to trigger the initial eruption by emerging flux.

Forbes (1989) has conducted a numerical experiment on emerging flux in which new bipolar flux is injected from below into a uniform horizontal field. Flux pile-up reconnection occurs with two streams of plasma produced by the reconnection and slowed down by termination shocks. The new flux rises and pinches off, forming a plasmoid and eventually approaching a potential field of lowest energy. The power output shows a slow increase in a driven phase, followed by an impulsive bursty release of energy and a slow decline, similar to what is observed in flares.

Some numerical experiments of the creation of hot flare loops have focussed on the close-down process by beginning with a line-tied open vertical field and watching it reconnect (Forbes and Priest, 1982; Forbes *et al*, 1989). Recent results show: an impulsive bursty regime which explains sudden jumps in loop height; a fast shock which slows the downflowing stream, increases the density and triggers a condensation; a reversed deflection current which deflects the downflow around a stagnation region; the splitting of the slow shock into a conduction front and an isothermal subshock; and the presence of evaporation in the hot loops to produce the observed density and temperature. Other numerical experiments (Mikic *et al*, 1988; Biskamp and Welter, 1989; Forbes, these proceedings) have modelled the global eruption followed by reconnection.

4. Conclusion

We have seen that reconnection may play many different roles on the Sun in a wide variety of phenomena from dynamo action to solar flares. At the present time new theoretical advances are being stimulated by surprises in numerical experiments, whose understanding is in turn being deepened by such theories. In future a better understanding is hoped both for the coupling to microscopic processes such as particle acceleration, for three-dimensional and time-dependent features and also for a closer comparison with observations.

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DISCUSSION

MOGILEVSKIJ: (i) The time of reconnection is usually short but in some flares the proton acceleration continues for longer than 5000 sec.
(ii) Can you give experimental arguments that the energy for the solar flare will be *only* by the reconnection of magnetic fields in the corona?

PRIEST: (i) Impulsive bursty reconnection in the impulsive phase may accelerate particles rapidly over short times, but then the quasi-steady reconnection in the main phase may continue to accelerate the particles over longer periods.
(ii) One cannot *prove* the existence of reconnection in a flare, since we shall not this century be able to resolve the reconnection region or map the field lines observationally with high precision. However, there are many indirect observations which are consistent with our theoretical understanding and with numerical experiments, and so reconnection is highly likely, especially since there is no other way of releasing the required energy from the magnetic field.

WEISS: As you know, the latest numerical experiments by Biskamp show complicated dynamical behaviour with self-similar structure associated with resistive instabilities. Biskamp asserts that reconnection has little resemblance to your models. Do you have any comments on this?

PRIEST: Biskamp's experiments are excellent but their interpretation is a matter of lively debate. Numerical and theoretical approaches complement one another and each has advantages and limitations. For example, the numerical experiments apply to a small range of magnetic Reynolds number and for specific boundary conditions. My view is that the theory does indeed enable us to understand the experiments better and, in particular, isolates the effect of a wide range of different boundary conditions. In particular, Forbes and Priest (1987, *Rev. Geophys* **25**,1583) have indeed been able to interpret the results of Biskamp in terms of their models. I would expect new reconnection experiments in future with a more comprehensive treatment of the boundary conditions to show Biskamp's results to be part of a wider class of solutions. In particular, we have shown how the type of reconnection depends crucially on the details of both inflow and outflow boundary conditions.

SMITH: (i) D. Biskamp is of the opinion that for 2-D steady reconnection you cannot exceed the Sweet-Parker rate.
(ii) Also your classification depends on a small parameter which the numerical experiments show is not small.

PRIEST: (i) What we have shown conclusively is that you can indeed exceed the Sweet-Parker rate provided you have the appropriate boundary conditions. Indeed you can obtain regimes similar to the Petschek regime or much faster regimes. It is not surprising that Biskamp did not find Petschek-like reconnection when you see what boundary conditions he has imposed. However, what I have presented here are, in a way, generalisations to the Petschek analysis to include extra effects that one finds in the experiments of Biskamp and others.
(ii) I agree that the Priest-Forbes theory does rely on assuming a small deviation from a uniform magnetic field. However, the resulting physical understanding in terms of slow- and fast-mode expansions is more general. The new Priest-Lee models are an attempt to drop the assumption of a nearly uniform field and we plan to try and generalise them even further.

FORBES: I have myself looked at Biskamp's numerical experiments in detail, and I believe they do show that you can get reconnection rates faster than the Sweet-Parker rate and the Petschek rate. Biskamp's definition of the reconnection rate is not the same as the standard one we use. When I re-evaluated his reconnection rate using the standard definition (Forbes and Priest, *Rev. Geophys*,**25**,1583 1987) I found that he does indeed have reconnection rates in excess of both the Sweet-Parker and Petschek rates in some of his cases.

LANG: (i) What is the driving force that makes magnetic fields come together?
(ii) Why does reconnection occur rarely?
(iii) Does reconnection happen in the photosphere, or in the corona?
(iv) Is there any observational evidence for magnetic fields actually coming together?

PRIEST: (i) It depends on the application. For example, in X-ray bright points the driver is the motion of photospheric magnetic footpoints, whereas in a two-ribbon flare I suggest that the eruption of a prominence drives reconnection below it.
(ii) I suspect that it actually occurs quite often, continually relaxing the field, reducing its complexity and energy and heating the corona.
(iii) In both. Photospheric magnetic fields tend to be dominated by plasma motions, which may drive the fields together, as in cancelling magnetic fields. Coronal magnetic fields tend to dominate the plasma and so reconnection can occur there either spontaneously by, say, the tearing mode instability or in response to footpoint motions.
(iv) Yes, for example in cancelling magnetic features and also there are cases where opposite polarity regions approach one another in active regions before flares.

KUNDU: I can cite several examples of evidence for magnetic field reconnection in solar flares, from VLA observations. The preflare region is generally heated (up to 10^7K) prior to flare onset; however, unless there is a newly emerging flux with appropriate polarity, the region does not flare. Here the new flux interacts with a pre-existing flux to produce a current sheet, which reconnects and produces a flare. Similarly, impulsive release of energy takes place in a gradual burst source when new flux emerges, produces a current sheet and leads to an impulsive burst via reconnection. These evidences are primarily dependent upon obtaining a spatial resolution of 1 arcsec.

UBEROI: There is a considerable literature on time-dependent magnetic reconnection in magnetospheric physics. Have these theories been applied to solar physics?

PRIEST: Yes, certainly. I have limited myself to steady reconnection but nonsteady effects have been considered analytically and numerically. The numerical studies of reconnection resulting from tearing or from an eruptive instability are strongly time-dependent. For regimes in which the central current sheet is small enough one tends to produce a quasi-steady reconnection, but when the central sheet is too large it goes unstable to secondary tearing (Forbes and Priest,1987), in which case one has an impulsive bursty modulation of the basic quasi-steady reconnection.

GOKHALE: (i) To what extent is the reconnection rate enhanced in your model?
(ii) If there are radiative losses from the reconnection region, will it not enhance the reconnection rate further?

PRIEST: (i) In a driven situation the reconnection may be slow or fast, depending on the speed of the driver, but for the almost-uniform models when $b < 0$ the maximum rate exceeds the Petschek value. For the nonuniform Priest-Lee models, for a given outflow speed the reconnection rate (*i.e.* inflow speed at large idistances) is a maximum at the smallest values of the central sheet length and it is typically a half of the outflow speed.
(ii) The energy balance will determine the plasma temperature in the central current sheet and so the magnetic diffusivity, but in a driven or freely reconnecting situation the main effect is to determine the size of the sheet rather than the rate of reconnection.

MONTGOMERY: We did a few preliminary 3D computations (incompressible MHD, periodic boundary conditions, initial-value problems) in order to watch the development of 3D perturbations of 2D current sheets. We saw very little that we could recognize from the well-studied 2D problem. (This involved no stabilizing, externally-applied, dc magnetic field.) The time-scales were fast, and the current sheets and field-surface perspective plots soon began to look like Swiss cheeses. What features of the 2D "X-point problem" do you expect to be able to recognize in 3D?

PRIEST: I agree very much that we need to go to 3D with numerical experiments and these will certainly give some surprises, just as the 2D experiments have done. Also, it is important to have a good interplay between numerical experiments and analytical models, using the former as a guide to constructing the latter if possible. However, the 2D calculations are still of great value since one can go to much higher magnetic Reynolds number and higher resolution and so they complement the 3D experiments. I would expect that locally near current sheets the 2D results may be valid in many cases. In some applications, such as coronal heating, the 3D aspects should be crucial, but in others such as two-ribbon solar flares with straight Hα ribbons and arcades the process is essentially 2D with a slow 3D modulation.

DRYER: Will reconnection occur only when an external forcing function or configuration is present?

PRIEST: No. In a sheared magnetic field or a current sheet, reconnection can be driven locally by a resistive instability. Indeed, the special feature of the Petschek regime is that it is, in a sense, "free" reconnection because it is the only regime where all the MHD characteristics are propagating information from the reconnection region.

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