HEATING THE SOLAR X-RAY CORONA

E.N. Parker

Dept. Physics, University of Chicago, Chicago, Illinois 60637, USA

ABSTRACT

The X-ray corona of the Sun consists of the diffuse X-ray background and the bright X-ray loops (10^7 erg/cm^{2}sec) confined in the strong (100 Gauss) bipolar fields of magnetic active regions. The bipolar fields are rooted in the solar granules which continually intermix the photospheric footpoints of the bipolar fields and progressively interlace the field lines. The intermixing is too slow to produce MHD waves. The magnetic field is never far from static equilibrium. The interlacing of the field lines, on scales of $100 - 1000$ km, produces magnetic stresses within the field that drive the field toward internal surfaces of tangential discontinuity (current sheets). This is the familiar rapid reconnection process, in which the field gradients are driven to ever increasing steepness, so that the slight electrical resistivity soon eats up the interlacing components of the field as fast as the magnetic stresses can steepen the gradients.

We suggest that this process is the principal heat source responsible for the X-ray corona of the Sun. It predicts that there are large numbers of nanoflares and picoflares throughout the X-ray loops. Katsukawa and Tsuneta find evidence in the Yohkoh X-ray observations of myriads of brightness fluctuations of the order of $10^{21} - 10^{23}$ ergs over times of 100 sec., opening the way for direct quantitative observational studies of the phenomenon.

Key words: Sun; X-ray; corona ; heating.

1. INTRODUCTION

The quiet corona of the Sun seems to be heated largely by the microflaring among the tiny magnetic features swept into the boundaries of the supergranule convective cells (Martin, 1984; Porter, et al. 1987; Porter and Moore, 1988). The estimated heat input to maintain the quiet corona is about 0.5×10^6 ergs/cm^{2}sec (Withbroe and Noyes, 1977). There appears to be sufficient energy in the microflaring for that purpose, although more precise observational studies are desirable. One imagines a suitable frequency spectrum of electromagnetic waves emitted from the microflares to provide heating both near and far out in the corona. The high ion temperature measured in the solar wind at 1 AU indicates active heating for some considerable distance out into space.

The recent analysis of white light corona pictures by R. Woo at the Jet Propulsion Laboratory adds a new dimension to the theory of heating in coronal holes. Comparing high resolution pictures of the corona taken at intervals of a few minutes, he finds strong variations in the individual radial striations, suggesting that the striations may be part of the coronal heating process. The characteristic time is about the same as the life of the photospheric granule. We note that an Alfvén wave created by a granule in a 10 Gauss field evolves, according to the WKB approximation, into a supersonic wave elongated radially as a consequence of the large Alfvén speed in the corona.

The active X-ray corona, on the other hand, confined within the bipolar magnetic fields (10^2 Gauss) of active regions, requires a heat input of the order of 10^7 ergs/cm^{2}sec (Withbroe and Noyes, 1977), i.e. about twenty times more than the quiet corona. So the microflaring is evidently not the major source of heat. The X-ray emission spectrum of the X-ray corona indicates that the temperature of the emitting gas is strongly inhomogeneous, ranging over $1 - 5 \times 10^6$ K across the fine (unresolved) filamentary structure along the field.

The photospheric convective cells - the granules - are the only known source of free energy at the solar surface with enough vigor to provide 10^7 ergs/cm^{2}sec. The characteristic scale l of the granules is 5×10^2 km, and the characteristic gas velocity v is of the order of 1 km/sec. Thus the characteristic granule turnover time $\tau = l/v$ is of the order of 5×10^2 sec. The gas density is $\rho \approx 2 \times 10^{-7}$ gm/cm^3, so the kinetic energy density is $\frac{1}{2}\rho v^2 \approx 10^3$ ergs/cm^3 and the characteristic convective energy flux is $\frac{1}{2}\rho v^3 \approx 10^8$ ergs/cm^{2}sec, or about 10 times the heat input to the active corona. Consequently the granules have been viewed as the likely source of heat responsible for the active X-ray corona. The question is how to convert the convective energy flux $\frac{1}{2}\rho v^3$ into heat in the million degree X-ray corona.

One possibility is Alfvén waves (cf. Alfvén, 1947) introduced into the 100 Gauss bipolar magnetic field by the convective motions at the photospheric footpoints of the field. So one approach has been to understand how an Alfvén wave, with a photospheric velocity amplitude comparable to the 1 km/sec granule velocity, might dissipate in the bipolar field of the X-ray corona. Wave damping by resistivity, anomalous resistivity, viscosity, and thermal conductivity have been considered, along with the possibility of resonance, phase mixing, etc. The fundamental difficulty with the concept is simply that the convective deformation of the bipolar field is too slow to generate Alfvén waves . The Alfvén transit time around a bipolar field of total length h is h/C, where C is the Alfvén speed $B/(4\pi\rho)^{1/2}$. The Alfvén speed in the X-ray corona is typically 2×10^3 km/sec (10^{10} hydrogen ions/cm^3, 10^2 Gauss), so the characteristic propagation time around a bipolar field of length $h = 10^5$ km is only 50 sec. Thus the slow moving, long lived granule, with a correlation time of the order of 500 sec, provides only a slow quasi-static deformation of the bipolar field. To put it another way, a wave with a period of 500 sec would have a wavelength of 10^6 km, well in excess of the 0.7×10^6 km radius of the Sun.

Now the granules are turbulent, with Reynolds numbers of the order of 10^8, so there must be smaller shorter lived eddies within each granule. Assuming the classical Kolmogoroff spectrum, in which the velocity $v(\lambda)$ of the eddies with characteristic scale λ is proportional to $\lambda^{1/3}$, it follows that the eddy life or correlation time $\tau(\lambda)$ is proportional to $\lambda^{2/3}$. The energy flux $\frac{1}{2}\rho v^3$ is proportional to λ, and therefore proportional to $\tau^{3/2}$. So a reduction in correlation time by a factor 10 reduces the power level by a factor of about 30, and the available energy becomes insufficient.

2. QUASI-STATIC FIELD DEFORMATION

It would appear that the basic quasi-static deformation of the bipolar magnetic field must, somehow, provide the principal heat input to the X-ray corona. The random shuffling and intermixing of the photospheric footpoints of the bipolar field continually interlace the field, providing the untidy field line topology sketched in Fig. 1 and the associated magnetic free energy. The interlacing takes place on the granule scale of 5×10^2 km. The question is how does the interlacing produce sufficient magnetic dissipation to supply the 10^7 ergs/cm^{2}sec to maintain the X-ray corona? The magnetic field at the photosphere is in the form of intense and widely spaced magnetic fibrils, with characteristic diameters of the order of 100 km and field intensities of the order of 1500 Gauss. Thus, in an active region where the mean field is 100 Gauss, the magnetic fibrils occupy a fraction 0.07 of the photosphere. So the fibril separation is about 400 km, or roughly the same as the characteristic scale, or correlation length, of the granules. Hence, we expect the individual fibrils to move with some degree of independence in the randomly shifting convective motions of the granules. The interlacing of the flux bundles in the bipolar magnetic field above the photosphere, sketched in Fig. 1, is the result.

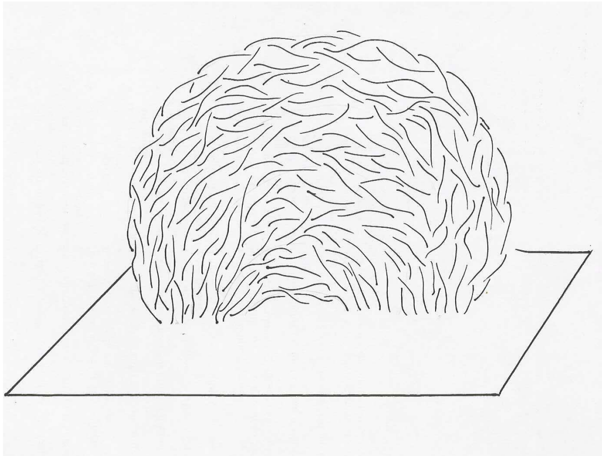

Figure 1. A schematic drawing of the magnetic field lines in the bipolar field of an active region, showing the interlacing produced by the random convective displacements of the photospheric footpoints of the field.

Now the coronal flux bundle defined by an individual fibril is displaced among the neighboring flux bundles, soon developing an inclination Θ relative to the mean field direction, as do the neighboring flux bundles as well. Then, if the individual flux bundles are not strongly twisted, they are each flattened by the pressure of their misaligned neighbors. We pointed out some years ago (Parker, 1981a,b) that the flattening would reduce the thickness γ of a bundle, thereby enhancing the resistive dissipation. For the record, then, the characteristic resistive dissipation time over a scale γ is $\gamma^2/4\eta$, where the resistive diffusion coefficient $\eta = c^2/4\pi\sigma$ in terms of the electrical conductivity σ. For fully ionized hydrogen $\eta \approx 0.5 \times 10^{13}/T^{3/2}$ cm^2/sec. Thus, at coronal temperatures $\eta \approx 10^4$ cm^2/sec. It follows that the resistive dissipation over 100 km is 3×10^9 sec, or 10^2 years. If γ is as small as 1 km, the dissipation time is 3×10^5 sec, or 3 days. We wondered if the flattening of the flux bundles might provide sufficient dissipation. In retrospect, it is not clear that the flattening and thinning of the individual flux bundle can go far enough.

For instance, the cross sectional area of the flux bundle is typically of the order of 10^4 km. Thus, if the thickness is reduced to 1 km, the width becomes 10^4 km, and it is not clear how so many thin sheets could be accommodated in the interlaced topology of Fig. 1, remembering that the over all scale of the bipolar region might not be much larger than 10^4 km. That is to say, wide sheets stack up like paper, rather than tangle and interlace like strings.

The concept of interlaced and mutually inclined flux bundles has been revisited by Priest, Hayvaerts, and Tile (2002) using up date observations of the magnetic fine structure, where they refer to the dislocation of the individual flux bundles as *flux tube tectonics*. They point out the additional complication that the magnetic flux of any given photospheric magnetic fibril at one end of a bipolar field may be presumed to connect into two or more fibrils at the other end. Since each of the connected fibrils moves independently, one expects that the flux bundle from a fibril splits directly into several flux bundles all heading for different destinations at the other end of the bipolar field (Longcope, 2001; Close, Hayvaerts, and Priest, 2004; Close, Parnell, and Priest, 2004).

The problem, then, is to estimate the rate of dissipation of magnetic energy at the current sheets. Passive resistive diffusion at the discontinuity where two nonparallel magnetic field components $\pm B$ press together produces an evolution of the field described by the error function $B\mathrm{erf}[x/(4\eta t)^{1/2}]$ if the plasma does not move, where x represents distance measured perpendicular to the surface of discontinuity at $x = 0$. The current density is

$$j(x,t) \;=\; \frac{cB}{4\pi(4\pi\eta t)^{1/2}} \exp\left(-\frac{x^2}{4\eta t}\right) \;. \qquad (1)$$

The dissipation rate is

$$\frac{j(x,t)^2}{\sigma} \;=\; \frac{B^2}{16\pi^2 t} \exp\left(-\frac{x^2}{2\eta t}\right) \;\text{ergs/cm}^3\text{sec}. \qquad (2)$$

The dissipation per unit area is

$$2\int_0^\infty dx \frac{j^2}{\sigma} \;=\; \frac{B^2}{8\pi}\left(\frac{\eta}{2\pi t}\right)^{1/2} \;\text{ergs/cm}^2\text{sec}, \qquad (3)$$

and the total energy dissipated in the time t is $(B^2/8\pi)(2\eta t/\pi)^{1/2}$ ergs/cm^2 . The essential point is that the rate of conversion of magnetic energy into heat declines as $1/t^{1/2}$ with the passage of time. Therefore, without some ongoing dynamical effect, e.g. unlimited flattening, to steepen the field gradient or rapid reconnection, the rate of dissipation quickly falls to negligible levels.

3. UBIQUITOUS RAPID RECONNECTION

As it turns out, there is a general dynamical effect introduced by the Maxwell stresses in the interlaced field line topology of the bipolar magnetic field of Fig. 1, so that the details of the interlacing become relatively unimportant. That is to say, the Maxwell stresses in the magnetic field automatically drive the field gradients to steepen without limit as a consequence of the expected interlaced field topology (Parker, 1972, 1983, 1988). The dissipation of magnetic energy does not decline to zero with the passage of time, but remains at some relatively high level. So the magnetic field can be everywhere continuous initially, and, if the field topology involves any significant interlacing, then the dynamical relaxation to static equilibrium (the lowest available energy state) involves the formation of surfaces of tangential discontinuity (TD's) in the field, i.e. current sheets of vanishing thickness and unbounded intensity in the ideal case of a perfectly conducting ambient fluid. This is, of course, the familiar phenomenon of *rapid reconnection* of the magnetic field, and the cause of magnetic flares on the Sun and other stars. It is unavoidable in an untidy field topology. So the X-ray corona of the Sun seems to be heated by the nanoflares arising in the current sheets in the small-scale interlacing of the bipolar magnetic fields, thanks to the peculiar properties of the Maxwell stresses in the interlaced field topology.

To show why and how the relaxation to the equilibrium described by Eq. (7) causes rapid reconnection, consider the idealized situation (Parker, 1986) in which a uniform magnetic field B extends in the z-direction through an ideal infinitely conducting fluid from an infinitely conducting end plate at $z = 0$ to an infinitely conducting end plate at $z = L$. At time $t = 0$ switch on the 2D incompressible fluid motion,

$$v_x \;=\; +kz\frac{\partial\Psi}{\partial y}\;,\; v_y \;=\; -kz\frac{\partial\Psi}{\partial x}\;,\; v_z \;=\; 0\,, \qquad (4)$$

where the stream function $\Psi = \Psi(x, z, kzt)$ is bounded, continuous, and n-times differentiable, but otherwise arbitrary. The endplate at $z = L$ participates in this motion. After a time t the magnetic field has the smooth bounded continuous form

$$B_x \;=\; +Bkt\frac{\partial\Psi}{\partial y}\;,\; B_y \;=\; -Bkt\frac{\partial\Psi}{\partial x}\;,\; B_z \;=\; B\,, \qquad (5)$$

sketched in Fig. 2. At time t we turn off the motion and hold both endplates $z = 0$, L fixed while releasing the fluid throughout $0 < z < L$ so that the field can relax to the lowest available energy state, i.e. stable equilibrium. A small viscosity is introduced to dissipate the motions over a period of time. The fluid pressure is maintained uniform at the endplates, with the result that the fluid pressure is uniform throughout $0 < z < L$.

In the final asymptotic equilibrium state the Maxwell stresses, described by the stress tensor

$$M_{ij} \;=\; -\,\delta_{ij}\frac{B^2}{8\pi} + \frac{B_iB_j}{4\pi}\;, \qquad (6)$$

are in equilibrium among themselves. The equilibrium state is described by

$$\nabla \times \mathbf{B}(\mathbf{r}) \;=\; \alpha(\mathbf{r})\mathbf{B}(\mathbf{r}) \;, \tag{7}$$

where the torsion coefficient $\alpha(\mathbf{r})$ is a scalar function of position.

Now this equilibrium equation is an unusual partial differential equation, having mixed characteristics. The curl of the equilibrium equation yields

$$\mathbf{B} \times \nabla\alpha \;=\; \nabla^{\mathbf{2}}\mathbf{B} + \alpha^{\mathbf{2}}\mathbf{B} \;, \tag{8}$$

while the divergence provides the simple restriction that

$$\mathbf{B} \cdot \nabla\alpha \;=\; \mathbf{0} \;. \tag{9}$$

It is evident from the Laplacian operator in Eq. (8) that the equilibrium equation has two families of complex characteristics. It is a quasilinear second order elliptic equation. On the other hand, it is evident from Eq. (9) that the torsion coefficient $\alpha(\mathbf{r})$ is constant along each field line, indicating that the field lines represent a family of real characteristics. We recall that a characteristic curve of a differential equation represents its own little world, with specification of the unknown function anywhere on a given characteristic curve determining the function everywhere else along the curve, but in no way determining the function on the neighboring characteristic curves. Thus the torsion coefficient α need not be continuous from one characteristic curve to the next.

So the equilibrium Eq. (7) has both real and complex characteristics, and therefore has properties beyond the more familiar equations of physics, which are usually either purely elliptic or purely hyperbolic. We must keep this in mind as we proceed with the investigation.

Now the solar granules introduce random swirling and mixing of the footpoints of the field, so we expect to find the field lines to be wrapped about their neighbors first one way and then the other at different locations along each field line, as suggested by Fig. 1 and Fig. 2. This is troubling because the torsion coefficient α is a measure of the local circulation of magnetic field around the direction of the field, which presumably varies in sign as a field line circulates in one direction and then the other around the neighboring field lines. However, Eq. (9) asserts that in equilibrium the torsion coefficient does not vary along the field. Yet with the field tied at both ends $z = 0, L$ and no resistive dissipation of the field, there is surely an equilibrium. This fact seems to bd contradicted by Eq. (9).

To be clear on the physical implications of the torsion coefficient, let Γ represent the magnetic circulation

$$\Gamma \;=\; \oint_C d\mathbf{s} \cdot \mathbf{B} \tag{10}$$

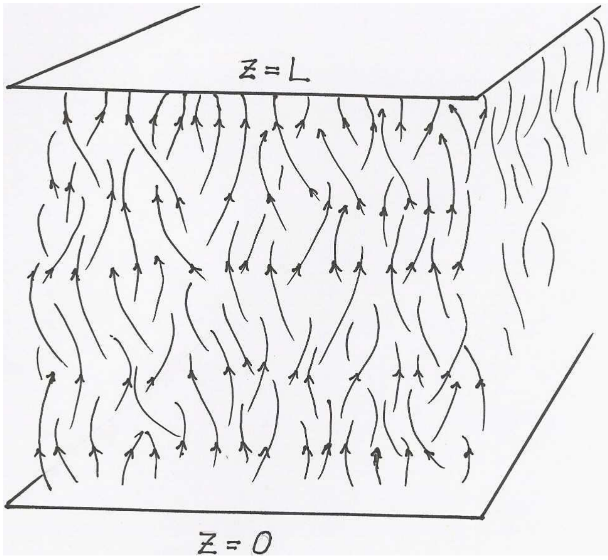

Figure 2. A schematic drawing of the interlaced field lines produced by the 2D incompressible motion described by the stream function $\Psi(x, y, kzt)$.

around a small closed contour C circling the field $\mathbf{B}$ at some point P. In the limit of a small contour and a continuous field $\mathbf{B}$, it follows from Stokes' theorem and Eq. (7) that $\alpha = \Gamma/\Phi$, where Φ is the magnetic flux across the surface S enclosed by the contour C

$$\Phi \;=\; \int_S d\mathbf{S} \cdot \mathbf{B} \;. \tag{11}$$

It follows that the torsion coefficient α represents the magnetic circulation per unit magnetic flux, Γ/Φ .

Obviously, if the field circles one way around the neighboring field at one location along the field and the other way somewhere else along the field, then α changes sign along the field in direct contradiction to Eq. (9). There can be no equilibrium, then, in a magnetic field exhibiting interlacing topology. To investigate this contradiction further, we dilate the system in the z-direction by the large factor Q ($\gg 1$). This has little effect on the z-component of the magnetic field, but it expands and diminishes the transverse field intensities B_x and B_y by the factor Q. The derivative $\partial/\partial z$ is diminished by the same factor, while the transverse derivatives $\partial/\partial x$ and $\partial/\partial y$ remain essentially unchanged. In the limit of large Q the equilibrium Eq. (9) reduces to the two dimensional vorticity equation for an ideal incompressible inviscid fluid, showing that α evolves with z in the same way that the vorticity ω evolves with time t. The vorticity equation has been the subject of intense mathematical investigation over the years (see review by Kraichnan and Montgomery, 1980), showing how the vorticity varies with time t, the enstrophy ω^2 migrating to ever smaller scales. Thus in the equilibrium magnetic field it follows that α^2

evolves with increasing z in the same way.

Now the interlacing of the magnetic field and the associated α are determined ahead of time by the arbitrary function $\Psi(x, y, kzt)$, which does not have the proper time evolution to provide the proper z variation for α to satisfy the vorticity-like equilibrium equation. Only if we precisely tailor the interlacing topology of the field to satisfy the special vorticity requirements can there be an equilibrium. It is evident, then, that the mathematical solutions to the equilibrium equation cannot be applied to the magnetic field created by the arbitrary successive swirls introduced by $\Psi(x, z, kzt)$. But physically we know that an equilibrium exists for a field tied at both endplates and preserved in the infinitely conducting fluid between.

Evidently we have introduced an unnecessary constraint into the mathematics, thereby excluding the possibility of a mathematical solution of the equilibrium Eq. (7). Recalling that the field lines represent a family of real characteristics and that specification of the torsion coefficient α on any one field line does not constrain α on the neighboring field lines, consider the possibility that the mathematical solution of the equilibrium equation contains surfaces of tangential discontinuity (TD's), i.e. shear planes and hence current sheets, across which the magnitude, but not the direction, of the field is continuous. Note, then, that a TD represents the surface of contact between the two regions of continuous field on either side. So neither the field nor its curl is defined on the TD. This releases us from the contradiction implied by Eq. (9) for a continuous field. Any variation in torsion or circulation along the field is accommodated by the unrestricted variation of the field shear across the TD.

We are obliged to conclude, then, that the relaxation of the interlaced magnetic field to equilibrium involves the creation of TD's by the Maxwell stresses in the field. Since we know from simple physical considerations that in the absence of resistivity the field relaxes asymptotically to a final equilibrium, it follows that the magnetic stresses must be of such form as to push the fluid to form the necessary TD's to create a final equilibrium. That is, of course, the basis for ongoing rapid reconnection and dissipation of magnetic field in the real world, with the Maxwell stresses continually sharpening the field gradients as the resistivity diffuses them. It follows that the expected interlacing of the field lines in the bipolar magnetic fields of active regions is an effective means for dissipating the free energy of the interlacing into heat in the X-ray corona. The formal mathematical treatment of the TD raises the question of the physics of the formation of the TD. The formation of a TD at a 2D X-type neutral point is sketched in Fig. 3. Increased pressure across the X-type neutral point (Fig. 3a) deforms the neutral point, sketched in Fig. 3b, into two Y-type neutral points with a TD extending between.

The problem can be treated formally using the optical analogy, which states that a field line in the equilibrium field described by Eq. (7) follows the same path through space as an optical ray path in a medium with an index

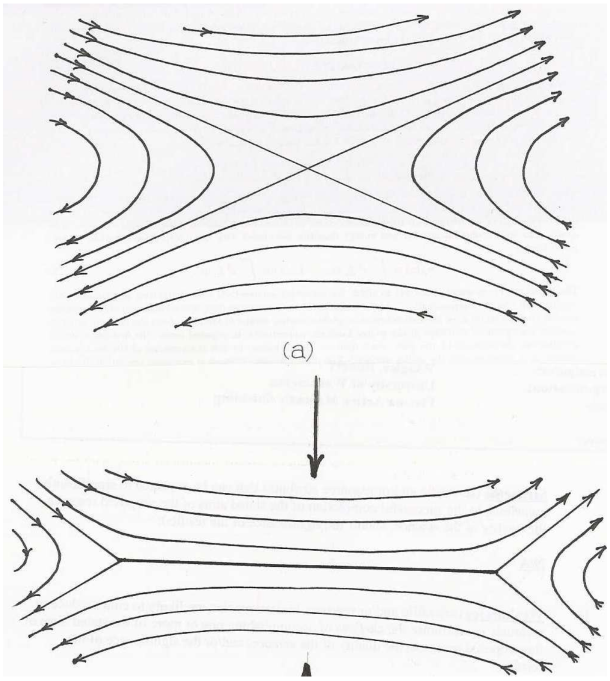

Figure 3. (a) A sketch of the field lines in the vicinity of the 2D X-type null point in a 3D magnetic field. (b) The flattening of the X-type null point by an enhanced local pressure, forming two Y-type null points with a TD between.

of refraction proportional to the field magnitude $B(\mathbf{r})$ (Parker, 1989a,b, 1990,1991, 1994). Thus the local maxima distributed throughout a field with untidy topology tend to exclude field lines, the optical path length being shorter if the field lines go around, rather than through, the local maximum (Fermat's principle). The region of avoidance creates a gap in the flux surfaces, allowing the continuous fields on either side to come into contact through the gap. Those two fields are generally not parallel where they meet in the gap, so their contact surface becomes a TD. That is illustrated in Fig. 3b, with the TD caused by the squashing of the initial X-type neutral point, i.e. the expulsion of field from between the two sectors of continuous field so that they no longer meet at a point as in Fig. 3a.

A relatively weak local field enhancement ΔB ($\ll B$) with scale w is sufficient to create an exclusion of field when the field lines are anchored at some large distance λ ($\gg w$) in either direction from the maximum. It is easy to show that the condition for excluding the field and creating a gap is

$$\frac{\Delta B}{B} > \frac{w}{\lambda} \tag{12}$$

in order of magnitude.

4. APPLICATION TO THE SOLAR CORONA

To illustrate how the creation of TD's in the untidy topology of a bipolar magnetic field of an active region of the Sun may heat the X-ray corona, consider the idealized example illustrated in Fig. 4. We show a single flux bundle attached to a wandering footpoint that moves with constant speed v (representing the granule motions) along a random meandering path among a static forest of fixed vertical flux tubes. The upper end of the wandering flux bundle is fixed at a height L. After a time t the wandering footpoint has traveled a distance vt and the flux bundle is inclined to the vertical by the angle Θ, where

$$\tan\Theta \;=\; \frac{vt}{L}\,. \tag{13}$$

The vertical field intensity is denoted by B, so that the inclination Θ is associated with the transverse (horizontal) component $B_\perp = B\tan\Theta$. The tension in the wandering flux bundle pulls back against the forward motion v of the footpoint with a force per unit horizontal area given by the Maxwell stress

$$F \;=\; \frac{B_\perp B}{4\pi}\,. \tag{14}$$

It follows that the forward motion v does work o the field at the rate

$$\begin{aligned} W \quad &= \quad vF\,, && (15)\\ &= \quad v\frac{B^2}{4\pi}\tan\Theta\,. && (16) \end{aligned}$$

In the Sun the individual flux bundles all resist the imposed convective swirling of their footpoints , so that the granules do work on the field at a rate of the order of our estimated W. The resistive dissipation at the incipient TD's in the bipolar field above the photosphere limits the increase in Θ, of course, and we can estimate the value of Θ from the necessity to do work W at the heating rate required by observations, viz. $W = 10^7$ ergs/cm^{2}sec. With $B = 10^2$ Gauss and $v = 1$ km/sec, it follows that $\tan\Theta = 0.1$ approximately, for which $\Theta = 6°$. This is quite a modest requirement. Unfortunately the scale of variation of Θ across the field is small, presumably of the order of 10^2 km or less, so direct observation is not possible at the present time.

Recent observations of the X-ray emission spectrum and the small-scale flickering in X-ray brightness appear to confirm the general picture, however. Both numerical simulations and laboratory experiments show that rapid reconnection tends to progress in bursts rather than smoothly. Thus we expect the X-ray corona to exhibit small scale flaring - nanoflares. To make a crude estimate of the energy release from an individual burst event

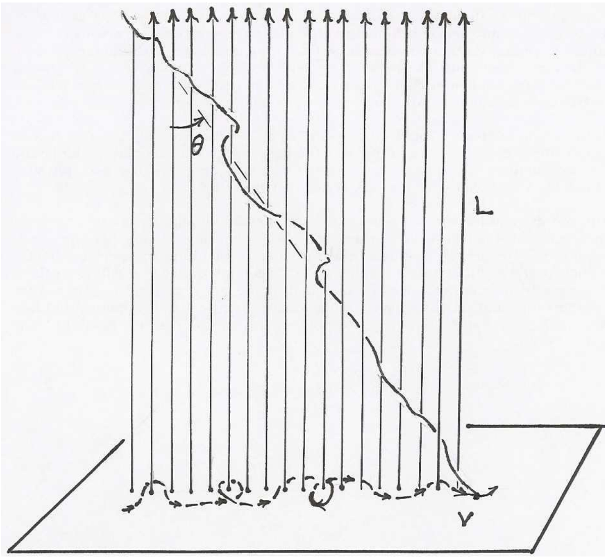

Figure 4. A schematic drawing of a single flux bundle whose photospheric footpoint wanders at random with velocity v among a forest of fixed vertical flux bundles. The upper end of the flux bundles are fixed at some large height L.

we estimate the magnetic free energy available around a local TD. Consider a volume with transverse scale l and extending a distance $l/\tan\Theta$ along the inclined local field. The associated volume is $l^3/\tan\Theta$ and the available magnetic energy density is $B_\perp^2/8\pi$. Thus with $\tan\Theta = 0.1$, $B_\perp = 10$ Gauss, and $l = 10^2$ km, the free energy is of the order of 4×10^{22} ergs. If a burst of reconnection were to dissipate one tenth of the available energy, the result would be a flaring event of 4×10^{21} ergs. So we might expect to observe bursts of energy over some range of the order of $10^{20} - 10^{24}$ ergs. The coronal temperature in the region would fluctuate up and down on some small transverse scale. The temperature rise would take place during the short life of the burst of reconnection - the nanoflare - and the subsequent cooling time at a density $N = 10^{10}$ /cm3 is estimated at perhaps 30 minutes.

It is interesting to note, then, that Katsukawa (2003; Katsukawa and Tsuneta, 2005) has studied the fluctuations in the individual pixels of the Yohkoh X-ray telescope. Katsukawa finds an excess over the expected thermal background noise, indicating rapidly varying X-ray emission with individual bursts in the range $10^{20} - 10^{24}$ ergs. Then some years ago, Sturrock, et al.(1990) and Feldman, et al (1992) analyzed the X-ray emission spectrum in some detail, finding that no combination of steady coronal temperatures could duplicate it. The problem was with the relative degrees of ionization. They pointed out that intermittent heating of the gas with short bursts of energy fitted the observed spectrum. So we seem to be looking at a sea of nanoflares, and, as Katsukawa points out, picoflares given that the standard large solar flare is 10^{32} ergs.

The general suggestion is that anywhere that magnetic fields become interlaced we may expect the formation of TD's with the associated rapid reconnection providing local heating and particle acceleration. Thus, for instance, the X-ray emission from other stars may be presumed to have an origin similar to the X-rays from the Sun. Then the active aurora seems to be associated with the onset of flux transfer effects in the terrestrial magnetosphere when the solar wind carries a south pointing magnetic field builds up the geotail with flux bundles stripped off the sunward magnetopause. We may reasonably suppose that the individual flux bundles are stretched out into the geotail in some degree of disarray, forming TD's whose lower ends we sometimes see as the auroral curtains (Parker, 1994). Other planets may be similarly affected.

5. REFERENCES

Alfvén, H. 1947, Granulation, magnetohydrodynamics, and the heating of the solar corona, Mon. Not, Roy. Astron. Soc. 107, 211-219.

Close, R. M., J. F. Hayvaerts, and E. R. Priest 2004, Effects of complexity on the flux-tube tectonics model, Solar Phys. 225, 267-292.

Close, R. M., C. E. Parnell, and E. R. Priest 2004, Separators in 3D quiet-Sun magnetic fields, Solar Phys. 225, 21-46.

Feldman, U., J, M. Laming, P. Mandelbaum, W. H. Goldstein, and A. Osterheld 1992, A burst model for the line emission in the solar atmosphere. II, Coronal extreme ultraviolet lines, Astrophys. J. 398, 692-697.

Katsukawa, Y. 2003, Spatial and temporal extent of solar nanoflares and their energy range, Pub. Astron. Soc. Japan 55, 1025-1031.

Katsukawa, Y. and S. Tsuneta 2005, Magnetic properties of footpoints of hot and cool loops, Astrophys. J. 621, 498-511.

Kraichnan, R. H. and D. Montgomery 1980, Two dimensional turbulence, Rep. Prog. Phys. 43, 547-619.

Longcope, D. W. 2001, Separator current sheets: Generic features in minimum-energy magnetic fields subject to flux constraints, Phys. Plasmas 8, 5277-5290.

Martin, S. F. 1984, in Small-Scale Dynamical Processes in Quiet Stellar Atmospheres, ed. S. L. Keil, Sacramento peak, Sunspot, New Mexico: National Solar Observatory.

Parker, E. N. 1972, Topological dissipation and the small-scale fields in turbulent gases, Astrophys. J. 174, 499-510.

Parker, E. N. 1981a, The dissipation of inhomogeneous magnetic fields and the problem of coronae. I Dislocation and flattening of flux tubes, Astrophys. J. 244, 631-643.

Parker, E. N. 1981b, The dissipation of inhomogeneous magnetic fields and the problem of coronae. II The dynamics of dislocated flux tubes, Astrophys. J. 244, 644-652.

Parker, E. N. 1983, Magnetic neutral sheets in evolving fields. II Formation of the solar corona, Astrophys. J. 264, 642-647.

Parker, E. N. 1986, Equilibrium of magnetic fields with arbitrary interweaving of the lines of force. I Discontinuities in the torsion, Geophys. Astrophys. Fluid Dynamics 34, 243-264.

Parker, E. N. 1988, Nanoflares and the solar X-ray corona, Astrophys. J 330, 474-479.

Parker, E. N. 1989a, Spontaneous tangential discontinuities and the optical analogy for static magnetic fields. II The optical analogy, Geophys. Astrophys. Fluid Dynamics 45, 169-182.

Parker, E. N. 1989b, Spontaneous tangential discontinuities and the optical analogy for static magnetic fields. III Zones of exclusion, Geophys. Astrophys. Fluid Dynamics 46, 105-133.

Parker, E. N. 1990, Spontaneous tangential discontinuities and the optical analogy in static magnetic fields. VII Effects of resistivity, Geophys. Astrophys. Fluid Dynamics 55, 161-182.

Parker, E. N. 1991, The optical analogy for vector fields, Phys. Fluids B 3, 2656-2659.

Parker, E. N. 1994, Spontaneous Current Sheets in Magnetic Fields, Oxford University Press, New York.

Porter, J. G., and R. L. Moore 1988, in Proc. 9th Sacramento Peak Summer Symposium, ed. R. C. Altrock, Sacramento Peak, Sunspot, New Mexco: National Solar Observatory.

Porter, J. G., R. L. Moore, E. J. Reichman, O. Engvold, and K. L. Harvey 1987, Microflares in the solar magnetic network, Astrophys. J. 323, 380-390.

Priest, E. R., J. F. Hayvaerts, and A. M. Title 2002, A flux tube model for solar coronal heating driven by the magnetic carpet, Astrophys. J. 576, 533-551.

Sturrock. P. A., W. W. Dixon, J. A. Klimchuk, and S. Antiochos, 1990, Episodic coronal heating, Astrophys. J. Lett. 356, L31-L34.

Withbroe, G. L., and R. W. Noyes 1977, Mass and energy flow in the solar chromosphere and corona, Ann. Rev. Astron. Astrophys. 15, 363-387.