MAGNETIC NEUTRAL SHEETS IN EVOLVING FIELDS. II. FORMATION OF THE SOLAR CORONA

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ABSTRACT

It is shown in the previous paper that whenever twisted flux tubes are bundled together, they are subject to dynamical nonequilibrium and internal neutral point reconnection, causing rapid dissipation of their torsion. In the present paper we explore the consequences of this general dynamical dissipation in the magnetic fields that produce the active corona of the Sun. The footpoints of the field are continually manipulated by the subphotospheric convection, so that the lines of force are continually wrapped and rotated about each other. The dynamical dissipation of the wrapping and rotation transfers the work done on the footpoints directly into heat in the corona, at a rate estimated to be of the order of $10^7$ ergs cm$^{-2}$ s$^{-1}$. The effect appears to be the principal source of heat to the visible corona.

This general picture implies that all magnetic fields extending outward from convecting astrophysical bodies produce intense heating of the tenuous outer atmospheres of those bodies, in general agreement with the observed fact of the universal activity of stars and galaxies.

*Subject headings: hydromagnetics — Sun: corona*

I. INTRODUCTION

It is interesting to reflect that three decades ago van de Hulst (1953) suggested that the key to coronal heating lies in the fine structure of the corona, and it is now almost 20 years since Gold (1964) suggested that coronal heating is, in some way, a consequence of the direct dissipation of the small-scale wrapping and twisting imposed upon the coronal magnetic fields by the convection in the photosphere. The problem has been that the high electrical conductivity of the corona frustrated any theoretical efforts to establish the necessary dissipation. Tucker (1973) explored some of the implications of Gold's idea, equating Joule dissipation of the field to the heat losses from the corona (radiation, conduction, and expansion). He pointed out that the dissipation of the magnetic field can supply the necessary energy to the corona only if the electric currents $\nabla \times \mathbf{B}$ are strongly concentrated, for some reason, into thin layers. The layers must be so thin, in fact, as to drive strong plasma turbulence to enhance the effective resistivity of the plasma. Tucker also gave a crude estimate of the rate at which the convective motions do work on the footpoints of the field at the photosphere. Levine (1974$b$) suggested that the fields develop internal neutral sheets somehow, which collapse and accelerate particles (Levine 1974$a$), thereby converting the energy of the magnetic strains into heat in the ambient coronal gas.

Rosner, et al. (1978) suggested that, in the presence of anomalous resistivity, the current must be compressed by a factor $10^5$ to $10^6$, from a characteristic scale of $10^8$ to $10^9$ cm to a thickness of the order of $5 \times 10^3$ cm. They make the point that, in addition to the continual introduction of small-scale strains into the field by the convection in the photosphere, the newly emerged and expanded flux tube is fed for a time from the torsion existing in the constricted tube beneath the surface of the Sun (Parker 1974, 1979, pp. 183–200). Rosner, Tucker, and Vaiana (1978) emphasize that detailed theoretical analysis of the observed structure of coronal loops, by themselves and others, allows no easy alternative to the idea that the visible corona is heated somehow by direct dissipation of the magnetic field.

More recently Sturrock and Uchida (1981) have attempted a quantitative estimate of the rate at which torsion is introduced in the flux tubes above the visible surface by the vorticity in the convective motions observed in the visible surface. They point out that if the flux tubes observed at the visible surface (with fields compressed to 1500 gauss in tubes with a radius of the order of 150 km) are rotated by the vorticity at peripheral velocities of 1 km s$^{-1}$ over a coherence time of $2 \times 10^3$ s (allowing about two full revolutions of the tube at the photosphere in each coherence time), then the magnetic energy input to the flux tube is equal to the observed radiative and conductive energy losses.

The mechanism causing the necessary concentration of the electric currents into suitably thin sheets to produce the desired dissipation has been addressed by

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Glencross (1975, 1980). He points out that the general stochastic shuffling of the footpoints of the flux tubes at the photosphere causes the tubes to be wrapped about each other in the corona in random patterns that vary along the length of the tubes. Hence, on the basis of the general nonequilibrium, and associated neutral point reconnection, of magnetic fields with a varying wrapping pattern (Parker 1972, 1979, pp. 359–391), he notes that the random wrapping of flux tubes about their neighbors is rapidly dissipated.

Subsequently, it has been pointed out (Parker 1981a, b) that the rapid dissipation of any individual flux tube that is misaligned relative to its neighbors by the stochastic displacements of its footpoints also contributes substantially to the heating of the corona.

The previous paper in this series (Parker 1982) makes the final point that even the torsion within aligned flux tubes is subject to nonequilibrium and neutral point reconnection with its neighbors. Altogether, then, it follows that all aspects of the small-scale strains introduced into the magnetic field by the subphotospheric convection are converted directly into heat in the corona, with the net result that the work done on the footpoints goes directly, with but little delay, into coronal heating. The dynamical nonequilibrium of close-packed, wrapped, and twisted flux tubes produces current sheets whose thickness declines asymptotically toward zero with the passage of time, until diffusion (presumably involving plasma microturbulence and other small-scale instabilities; see Priest 1981b) takes over and destroys the transverse components of the field. If the photospheric motions were to cease, the field would soon relax into a large-scale configuration with neither small-scale twisting nor wrapping of flux tubes.

When the case, there remains the basic question of the rate at which the subphotospheric convection introduces strains into the field by twisting and wrapping the individual flux tubes. Associating the twisting and wrapping with the granules suggests a basic transverse scale of the order of 10³ km at the photosphere. A similar scale arises from the observational picture, in which the individual flux tubes of 1500 gauss and 200–300 km diameter (Livingston and Harvey 1969; Stenflo 1973; Chapman 1973) expand to 10³ km above the photosphere to fill the available space with a mean field of, say, 10² gauss.

II. CORONAL HEATING REQUIREMENTS

The observable corona (both the quiet and the active regions) is dominated and confined by re-entrant magnetic fields extending up from the surface of the Sun. The magnetic fields cut off both outward thermal conduction and outward thermal expansion into space, which would otherwise drain away so much of the energy that the temperature and density could not build up to observable levels (the coronal hole is an example of the conditions that prevail in the absence of magnetic confinement).

The most intense coronal emission is from the regions of strong closed fields above active regions on the visible surface of the Sun. The coronal EUV and X-ray emission there may be 10–100 times more intense than from the closed quiet region, which in turn may be 10 times brighter than the coronal hole. Withbroe and Noyes (1977) provide estimates of the energy losses, summarized here in Table 1.

Observations show that the emitting gases are contained in fine striations or loops extending along the lines of force of the enclosing field. Each separate strand of radiating gas is isolated by the magnetic field from its neighbors, continually shifting relative to its neighbors (Rosner et al. 1978) and continually variable and active (Withbroe 1981; Priest 1981a). Rosner, Tucker, and Vaiana (1978) and Rosner et al. (1978) point out that the visible corona of the Sun is evidently a magnetic phenomenon made up of close-packed bundles of magnetically controlled, magnetically heated, and magnetically isolated filaments of relatively dense superheated gas.

Consider, then, an active region with fields of the order of \( B = 10^2 \) gauss extending a distance \( L = 10^{10} \) cm up into the corona. The large-scale field is composed of small-scale, close-packed, twisted, and mutually wrapped flux tubes. The magnetic energy per unit area of solar surface is of the order of

\[
W = \frac{LB^2}{8\pi} = 4 \times 10^{12} \text{ ergs cm}^{-2}.
\]

The active region is heated at a rate (see Table 1) \( H \approx 10^7 \) ergs cm\(^{-2}\) s\(^{-1}\). Hence, the magnetic energy \( W \), or its equivalent, must be dissipated in a characteristic time

\[
t_D = W/H = 4 \times 10^5 \text{ s}.
\]

It is only the strains from an equilibrium configuration that are available for dissipation—the magnetic free energy—which might be a quarter of \( W \) as defined above. Hence, we suggest that the characteristic dissipation time is more like \( t_D = 10^5 \) s, or less.

III. GENERATION OF MAGNETIC STRAIN

Consider, then, the fundamental question of the rate at which strains are introduced into the field by the displacement of the feet of the lines of force at the photosphere (see also the discussions in Tucker 1973 and Levine 1974a, b). If the production and dissipation

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of strains is to heat the active corona, requiring \(10^{7}\) ergs cm\(^{-2}\) s\(^{-1}\) (see Table 1), then the displacement of the footpoints must do work on the field at a rate of \(10^{7}\) ergs cm\(^{-2}\) s\(^{-1}\).

Observations deal with the so-called "magnetic knots" or "magnetic bright points" (Beckers and Schröter 1968), which we suppose to be the footpoints of the magnetic flux tubes filling the corona. Smithson (1973) made a detailed study of the motion of the knots in quiet regions. If we assume that the magnetic knots move about as much in active as in quiet regions, his results provide an estimate of the strain rate in active regions. Smithson reports (a) that the individual magnetic bright point jiggles about with a correlation time \(\tau\) small compared to 1 day \((\tau \approx 10^5\) s) and (b) that the individual knot undergoes large random displacements of \(5-20 \times 10^3\) km at intervals of a day or two. To be more specific, Smithson studied the bright vertices of the calcium network associated with the cluster of flux tubes extending through the surface of the Sun at the junctions of the supergranule boundaries. Smithson found random motions of the vertices (i.e., of the supergranule boundary junctions) equivalent to an rms displacement of \(8 \times 10^3\) km in 24 hr, equivalent to a turbulent diffusion coefficient \(K = 4 \times 10^{12}\) cm\(^2\) s\(^{-1}\). The sudden large displacements contribute an equal or greater diffusion. The foot of a flux tube is displaced an rms distance of at least \(10^4\) km in a 24 hr period.

During the 24 hr period the individual footpoints wander about among neighboring footpoints along a random path that is considerably longer than \(10^4\) km, of course. To be more precise, a random walk with steps of length \(\lambda\) at time intervals \(\tau\) produces an effective diffusion coefficient \(K \approx \frac{\lambda^2}{\tau}\). In a time \(t\) the rms displacement is \((4Kt)^{1/2} = (4\lambda^2t/3\tau)^{1/2}\), while the total path length traversed is \(\Lambda = \lambda t/\tau\). The total path length is larger than the rms displacement by the factor \((3t/4\pi)^{1/2}\). Thus, if the small random steps are at intervals of, say, 1 hr, then the total path is approximately a factor of 5 longer. An rms displacement of \(8 \times 10^3\) km in 24 hr suggests, then, a total random displacement of \(\Lambda = 4 \times 10^4\) km in the same period, as a consequence of the continual small-scale "jiggling" of the magnetic knots (in addition to the occasional large-scale displacements). The mean velocity is, accordingly, \(w = \lambda/\tau = 0.4\) km s\(^{-1}\). Figure 1 is a sketch of a single tube whose footpoint (at \(z = 0\)) has been randomly walked about for a time.

If the strain produced by this random motion of individual magnetic knots accumulates for a period \(t = 24\) hr, spreading out a distance \(L = 10^{10}\) cm along the length of the coronal flux tubes, then the lines of force are displaced by an angle \(\theta\) such that \(\tan \theta = \Lambda/L \approx 0.4\). The transverse component of the field is

\[ B_t = B \tan \theta = BwL = 0.4B, \]

and, together with the inhomogeneities introduced into the original longitudinal component \(B\), raises the magnetic energy by about one-fourth. It was pointed out in § II that coronal heating in a field of \(B = 10^2\) gauss above an active region requires the equivalent of the dissipation of one-fourth of \(B^2/8\pi\) in a 24 hr period. Hence, Smithson's observations of the motion of magnetic knots in quiet regions suggest a strain rate that is adequate for heating the corona. More detailed observations can only uncover further motions, involving rotation of the individual magnetic knots, etc., not included in Smithson's first cut at the problem (see Sturrock and Uchida 1981).
The calculation of the energy input in the corona can also be carried out in terms of the magnetic stresses at the photosphere where the fluid motions are driving the field. Consider a field $B$ extending vertically from the photosphere at $z = 0$ to where it is fixed at $z = L$, as sketched in Figure 1. It is evident that a net horizontal displacement of length $wt$ (along some random path) of the foot of a tube of length $L$ and strength $B$ produces an oblique tube of length $(L^2 + w^2 t^2)^{1/2}$ involving a mean transverse flux density $B_t = B wt/L$, as sketched in Figure 1. The magnetic stress $B_t B/4\pi$ opposes the horizontal displacement velocity $w$, so that the rate at which the displacement does work on the tube is

$$P = wB_t B/4\pi = B_t^2 w^2 t/4\pi L \text{ ergs cm}^{-2} \text{s}^{-1}.$$  

If dissipation is sufficiently slow that the reconnection does not begin to destroy $B_t$ until $B_t$ has accumulated for 1 day, then the numbers $B = 10^2$ gauss, $w = 0.4$ km s$^{-1}$, $L = 10^{10}$ cm already employed yield $P = 10^7$ ergs cm$^{-2}$ s$^{-1}$.

IV. RECONNECTION RATE

It was shown in the previous paper (Parker 1982) that the strain is dissipated rapidly as a consequence of the associated dynamical nonequilibrium, taking the form of magnetic neutral sheets at which neutral point reconnection occurs.

We are dealing with the tenuous coronal gas whose pressure $p$ is small compared to the pressure $B^2/8\pi$ of the magnetic field in which the gas is embedded. It follows that the characteristic Alfvén speed $V_A = B/(4\pi p)^{1/2}$ is large compared to the speed of sound and hence large compared to any of the motions that manipulate the field. Specifically, in an active region of the corona, where $N = 2 \times 10^9$ H atoms cm$^{-3}$, $T = 2 \times 10^6$ K, and $B = 10^2$ gauss, the Alfvén speed is 5000 km s$^{-1}$ and the sound speed is 240 km s$^{-1}$. The gas pressure is 1 dyne cm$^{-2}$, while the magnetic pressure is 400 dynes cm$^{-2}$. It was suggested in the previous section that the feet of the flux tubes may be manipulated at speeds of the order of 0.4 km s$^{-1}$. Such slow manipulation provides a slow evolution of the static equilibrium state of the field, with growing magnetic neutral sheets throughout. The essential feature of the magnetic neutral sheet is the expulsion of fluid from the sheet, causing the thickness $l$ to decline asymptotically to zero, progressively concentrating the electric current until dissipation becomes important. The dynamical theory of magnetic neutral sheets is well known (see reviews by Vasyliunas 1975; Parker 1979, pp. 359–436; Priest 1981b; Syrovatsky 1981). The essential point is that under quasi-steady conditions the thickness $l$ of the neutral sheet formed between fields of characteristic scale $h$ declines to a value somewhere in the interval

$$N_{R}^{1/2} \gtrsim 1/h \gtrsim N_{R}^{-1} \ln N_{R},$$

where $N_{R}$ is the magnetic Reynolds number $hV_A/\eta$ in terms of the effective resistive diffusion coefficient $\eta$. As
a consequence of this small value of $l$ the field moves into, and reconnects across, the neutral sheet at a speed $v = \eta/l$ somewhere in the interval

$$V_A/\eta N^2 R \lesssim v \lesssim V_A/\eta N_R.$$  

(1)

The precise values of $l$ and $v$ are determined by the peripheral conditions on the fluid pressure, etc. Note that with $h = 10^9$ cm, $V_A = 5000$ km s$^{-1}$, and $\eta = 1.4 \times 10^3$ cm$^2$ s$^{-1}$ for ionized hydrogen at $2 \times 10^6$ K, the magnetic Reynolds number has the large value $N_R = 4 \times 10^{14}$, so that equation (1) provides the broad range $5 \times 10^{-8} \lesssim v/V_A \lesssim 3 \times 10^{-2}$.

Suppose, then, that magnetic reconnection proceeds at some rate $v$ across the neutral sheets in $B_i$. The reconnection cuts across the diameter $h$ of a flux tube in a time $t = h/v$, limiting the growth of $B_i$ to

$$B_i = B_i(t/L) = B(h/L) w/v.$$

The rate at which work is done by $w$ is, accordingly,

$$P = w (B^2 / 4 \pi) (w/v) h/L.$$

With $h/L = 10^{-1}$ ($h = 10^9$ cm, $L = 10^{10}$ cm), $B = 10^2$ gauss, and $w = 0.4$ km s$^{-1}$, the requirement that $P = 10^9$ ergs cm$^{-2}$ s$^{-1}$ yields $v = 1.3 \times 10^4$ cm s$^{-1}$, with $B_i/B = 0.3$, implying a pitch angle $\theta = \tan^{-1} B_i/B = 20^\circ$. This reconnection rate seems modest. The lower and upper bounds on $v$, indicated by equation (1) for the estimated values $N_R \approx 4 \times 10^{14}$ and $V_A = 5 \times 10^8$ cm s$^{-1}$, are 25 and 1.6 $\times 10^7$ cm s$^{-1}$, respectively. The harmonic mean of these extremes is $2 \times 10^4$ cm s$^{-1}$, to be compared with our estimated requirement $v = 1.3 \times 10^4$ cm s$^{-1}$. Indeed, the point should not be overlooked that power input is inversely proportional to $v$, so that if $v$ were smaller, the power input would be larger! It would seem from these numbers that the observed coronal heating is a natural consequence of the rotation and random displacement of the flux tubes in the photosphere.

The calculations have been carried out for the active corona, in a region of $10^2$ gauss. The quiet corona follows automatically for weaker fields, of say 10 gauss, subject to the same manipulation at the photosphere. The energy requirement (see Table 1) is $\approx 1/30$ of the active corona.

There is an important point that should be noted here in connection with the production of the field in the field above the visible surface of the Sun. The strain rate has been estimated from observations of the random walk of individual magnetic knots, neglecting entirely the strains that may result from rotation of the individual flux tubes where they penetrate the photosphere. It has been pointed out elsewhere (Parker 1974, 1975, 1979, pp. 167–206) that any twisting of an intense flux tube beneath the surface of the Sun is quickly transported (as a torsional Alfven wave) to the expanded portion of the tube above the surface (Rosner et al. 1978). It follows, then, that coronal heating may be as much (or more) a consequence of turbulent effects beneath the visible surface of the Sun as it is a consequence of the random walk of the magnetic knots at the surface. Perhaps the higher resolution (of 0.1" $\rightarrow$ 75 km) of the planned Solar Optical Telescope can one day shed some light on this basic question.

There is an immediate prediction from this theoretical picture of coronal heating. The magnetic neutral sheets are formed as a consequence of the fluid being squeezed out of the sheets along the magnetic lines of force, so that the thickness $l$ declines asymptotically to zero. The ejection of fluid may proceed at some significant fraction of the speed of sound (Parker 1957). Hence, we expect small-scale high speed sheets of gas associated with the visible corona. Brueckner (1981) pointed out that observations at high resolution indicate the presence of small-scale, high speed fluid motions in the corona.

The principles are general, of course, so that we are led to the conclusion that there is strong heating of the outer atmosphere of any rotating convective object. Rotation and convection appear to be a sufficient condition for producing a magnetic field in any astronomical body (see discussion in Moffatt 1978; Parker 1979, pp. 710–712; Krause and Rädler 1980). Magnetic buoyancy guarantees that the fields extend upward through the surface of the body into the tenuous outer atmosphere (Parker 1979, pp. 314–355). Hence, any convective stirring of the field at the surface (or below the surface) introduces an accumulating strain into the extended field above. The automatic nonequilibrium of the extended field produces neutral point reconnection that transfers the magnetic free energy in the strains to heat in the tenuous gas.

This general theoretical assertion is in accord with observations, which show, for instance, the extensive neutral sheets and shear planes (current sheets) in active planetary magnetospheres as a consequence of strains introduced by the solar wind which drives the magnetospheric convection (Parker 1958; Axford and Hines 1962; Akasofu and Chapman 1972; Johnson 1978; Smith et al. 1980; Behannon, Burlaga, and Ness 1981; Sonnerup et al. 1981; Akasofu 1981). Coronal X-ray emission from a variety of stars indicates extensive heating, often far beyond the levels to be found in the Sun (see Bonnet and Dupree 1981). Hence, we suppose that the same magnetic effects are at work in other stars, in many cases in a more vigorous state.

Magnetic reconnection evidently occurs to some degree throughout the solar wind where the turbulence and the general nonuniformity of the wind produce severe strains in the field (Parker 1963; Hundhausen 1972,
REFERENCES


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