STRUCTURE OF CORONAL NEUTRAL SHEETS

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ABSTRACT

A qualitative model for the structure of the neutral sheet lying along the axis of coronal streamers is developed. The internal topology of the sheet is that of extremely thin magnetic tongues greatly distended outward by the solar wind expansion inside the sheet. Due to finite conductivity effects, expansion is taking place across the field lines but is retarded relative to the external flow by the reverse **j**X**B** force. The sheet thickness is determined by three considerations: the electrical conductivity that specifies the magnitude of the gradients in field strength, the expansion velocity that stretches the field lines outward decreasing the sheet thickness, and finally, the lateral pressure balance that limits the approach of the oppositely directed external field toward the neutral plane.

If σ is the electrical conductivity, the sheet thickness is shown to be proportional to σ-1/3. For an electron conductivity evaluated perpendicular to the internal field in the sheet, the thickness is of the order of 100 km in the inner corona and 10,000 km at 1 AU. Microturbulence and instabilities are expected to yield dimensions greater than these theoretical values since these effects tend to reduce the "effective" conductivity.

INTRODUCTION

It is clear that a great number of neutral sheets should exist in the solar corona, considering the numerous polarity reversals in the observed photospheric magnetic field. The intersection of these sheets with the ecliptic at 1 AU can be observed with satellites and are normally referred to as "sector boundaries." However, the detection of structures as thin as these in the inner corona would be difficult through normal photographic techniques. Since an eclipse photograph of the inner corona reflects an integration of electron density over the line of sight, these sheets, even though denser than their surroundings, would be essentially transparent unless they were not only planar but also had their plane coinciding excactly with the line of sight direction.

There seems to be an increasing amount of indirect observational evidence that such sheets may be the tracks along which high energy particles released during flares travel from the sun to the earth. For example, the most favored theory of type III radio bursts proposes that this emission is produced by a cloud of outward traveling electrons exciting plasma oscillations at successively higher levels in the corona. It has been proposed [*Wild,* 1964] that these electrons move out along the neutral sheets situated along coronal streamers. In addition, *Bumba and Obridko* [1969] have suggested that proton flare activity associated with Bartels' active longitudes occurs in the neighborhood of the sector boundaries of the interplanetary magnetic field.

In this paper, we describe a physical model of the expected structure of such coronal neutral sheets. Though the approximations are crude, it is hoped that the model will provide physical insight as to the grossest features that would be expected to be associated with these structures.

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DESCRIPTION OF THE MODEL

Since the structure of a neutral sheet essentially is produced by finite conductivity effects, the simplest way to visualize it is to first consider a coronal streamer for the case $\sigma = \infty$, introduce a small amount of resistivity to the plasma, and examine the consequences. A typical helmet streamer (fig. 1) consists of a region of closed magnetic loops in the corona with open field lines adjacent to and above the loops, and is commonly observed during solar eclipses. For $\sigma = \infty$, a neutral point exists at the top of the closed loops, and true sheet currents of zero

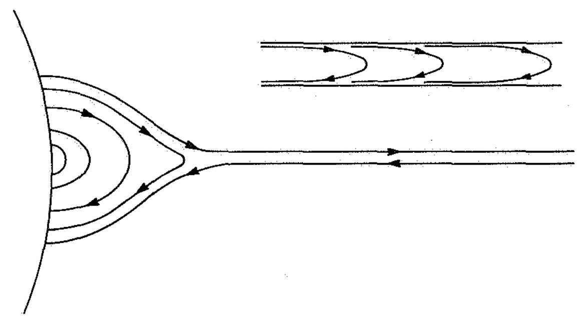

Figure 1. *Schematic of typical helmet streamers. The insert shows a blown up section of the neutral sheet along the steamer axis.*

dimension exist both above the neutral point and below it between the open and closed regions [*Pneuman and Kopp,* 1971]. These lower sheet currents exist because expansion occurs along the open field lines but not along the closed ones, resulting in a pressure discontinuity between the two regions that must be balanced by a jump in magnetic field strength. The location of the neutral point is determined by the condition that the pressure differential between the closed and open regions at that point is exactly balanced by the magnetic pressure in the open region; that is,

$$p_{cl} - p_{st} = \frac{B_{st}^2}{8\pi} \tag{1}$$

where *cl* refers to the closed region and *st* to the streaming region.

We now investigate the consequences of introducing a small resistivity into this model. Obviously, the first will be reconnection of open field lines at the neutral point. The outward part of the reconnected field will just be expelled outward along the sheet, whereas the inward part will become a new closed loop attached to the solar surface. Consequently, a hydrostatic pressure distribution will be set up along this loop. However, the new loop will lie outside that for which equation (1) is satisfied, and the pressure differential will be larger than the external magnetic pressure. As a result, the loop will expand outward along the sheet and, were the conductivity infinite, it would continue expanding to infinity (since the field and material are frozen). Hence, this loop would eventually become a new *open* field line and the configuration would revert to that which existed before reconnection took place. If, on the other hand, σ is finite, a small amount of relative motion between the field and gas is permitted and the loop will not be expelled to infinity but will only be pulled outward to some location where the diffusion of field backward through the material is just balanced by the outward convection of the field lines. When a steady state is achieved, the neutral sheet will have a finite thickness and will contain loops of weak magnetic field attached to the solar surface with the plasma expanding outward across the loops. The resulting configuration might look something like that shown at the top of figure 1. Since the relative diffusion of the field through the gas is expected to be very small for coronal conditions, these internal loops would be extremely distended and, consequently, the transverse field there should be very small.

If this physical picture of the neutral sheet is adopted, the sheet thickness will depend on three factors:

1. The magnitude of the electrical conductivity. This determines how large the gradients in magnetic field strength can be within the sheet. As one would expect, the sheet dimension should vary inversely with the conductivity.
2. The expansion velocity within the sheet. Large velocities will stretch the field lines more and decrease the sheet thickness.
3. The lateral pressure balance between the sheet and its surroundings. This consideration will limit the approach of the external field towards the neutral plane.

In the next section, we employ these concepts to develop a first-order model of coronal neutral sheets. We employ the MHD approach throughout. Since these structures are expected to be thin, one must ascertain whether collisions will be frequent enough to make this approximation valid. However, since the sheet dimensions are not known *a priori*, it seems most fruitful to investigate this question in the light of the properties of the resulting solution.

MAGNETIC FIELD CONFIGURATION

Consider a coronal streamer, situated at the solar equator, which is axisymmetric with respect to the rotational axis. If θ is the polar angle; then, the neutral sheet will lie along the plane $\theta = \pi/2$. As a first approximation, the sheet is assumed radial in cross section and the fluid velocity inside is also taken to be radial.[1] In line with the arguments of the previous section, the magnetic field

[1] *Throughout this work we omit the effects of solar rotation.*

will have both radial and transverse components. Hence, the velocity and magnetic field vectors are given respectively by

$$\mathbf{V} = V_r\hat{e}_r \quad \text{and} \quad \mathbf{B} = B_r\hat{e}_r + B_\theta\hat{e}_\theta$$

where $\hat{e}_r$ and $\hat{e}_\theta$ are unit vectors in the radial and polar directions.

For steady state conditions, the expression for conservation of magnetic flux and the induction equation are

$$\nabla \cdot \mathbf{B} = 0 \qquad (2)$$

$$\nabla \times \left[\mathbf{V} \times \mathbf{B} \cdot \left(\frac{1}{4\pi\sigma}\right) \nabla \times \mathbf{B}\right] = 0 \qquad (3)$$

where σ is the electrical conductivity. In deriving equation (3), we have taken σ to be a scalar. This would not be the case in the presence of a strong magnetic field, but the field inside the sheet is expected to be weak and the reduction in conductivity due to the field may not be great. This question will be considered again later so that actual conditions may be more realistically assessed. We merely consider here that the conductivity is a function of temperature only; we also assume isothermal conditions such that σ is a constant.

Subject to the above conditions, equations (2) and (3) can be solved in the limit $R_{em} >> 1$ to yield expressions for both the radial and transverse components of the magnetic field. We merely tabulate the results here.

$$B_r = \frac{B_o}{x^2} \exp\left[-\frac{n(n+1)}{R_{em}} \int_1^x \frac{dx}{x^2 u}\right] \frac{P_n(\cos\theta)}{P_n(\cos\theta_b)} \qquad (4)$$

$$B_\theta = \frac{-B_o}{R_{em} x^3 u} \exp\left[-\frac{n(n+1)}{R_{em}} \int_1^x \frac{dx}{x^2 u}\right] \frac{P_n'(\cos\theta)}{P_n(\cos\theta_b)} \qquad (5)$$

Here, $x = r/r_o$, where r_o is the radial distance to the top of the helmet in the streamer. This is where the sheet is assumed to originate and is where the neutral point would be located if the conductivity were infinite. B_o is the magnetic field strength at r_o just outside the sheet: $u = V_r/V_s$ (V_s being the sound speed); R_{em} is a magnetic Reynolds number given by $R_{em} = 4\pi\sigma r_o V_s$; $P_n(\cos\theta)$ is a Legendre polynomial of order n; and θ_b is the angle representing the sheet boundary. The prime denotes differentiation with respect to θ. (Note that for a coronal temperature of 1.0×10^6 °K and $r_o = 2R_\odot$, $R_{em} = 4.5\times10^{14}$. This is an enormous number and shows just how effectively the field and material are frozen in the corona.) To satisfy the two conditions that $B_r = 0$ for $\theta = \pi/2$ and $B_\theta = 0$ for $\theta = \theta_b$, we require that n be odd and also a root of the equation

$$P_n'(\cos\theta_b) = 0 \qquad (6)$$

Since n characterizes the variation of quantities across the sheet, it will be a large number inversely proportional to the sheet thickness. For example, if $n >> 1$, we can write

$$P_n(\cos\theta) \approx \left(\frac{2}{n\pi\sin\theta}\right)^{1/2} \sin\left[\left(n + \frac{1}{2}\right)\theta + \frac{\pi}{4}\right]$$

Condition (6) then reduces to

$$\cos\left[\left(n + \frac{1}{2}\right)\theta_b + \frac{\pi}{4}\right] = 0 \quad \text{or} \quad n = \frac{\theta_b + (\pi/2)}{\pi - 2\theta_b}$$

and since $\pi/2 - \theta_b = \delta/r$ where δ is the sheet thickness, we have

$$\delta \approx \frac{\pi r}{2n} \qquad (7)$$

CURRENT DENSITY, PRESSURE DISTRIBUTION, EXPANSION VELOCITY, AND SHEET THICKNESS

Having derived the magnetic field configuration, we can now derive the current density J from Ampere's law.

$$J = \frac{-B_o}{4\pi r_o x^3}\left\{1 - \frac{1}{R_{em} x^2 u^2}\left[\frac{d}{dx}(x^2 u) + \frac{n(n+1)}{R_{em}}\right]\right\}$$

$$\exp\left[-\frac{n(n+1)}{R_{em}} \int_1^x \frac{dx}{x^2 u}\right] \frac{P_n'(\cos\theta)}{P_n(\cos\theta_b)} \qquad (8)$$

Considering the lateral balance of forces across the sheet, we require that

$$\frac{1}{r}\frac{\partial p}{\partial\theta} = (\mathbf{J}\times\mathbf{B})_\theta \qquad (9)$$

where p is the gas pressure. Integration of equation (9) then yields an expression for the gas pressure in the sheet.

$$p = p_e(x) + \frac{B_o^2}{8\pi x^4} \exp\left[-\frac{2n(n+1)}{R_{em}} \int_1^x \frac{dx}{x^2 u}\right]$$

$$\left\{1 - \frac{1}{R_{em} x^2 u^2}\left[\frac{d}{dx}(x^2 u) + \frac{n(n+1)}{R_{em}}\right]\right\}$$

$$\left\{1 - \left[\frac{P_n(\cos\theta)}{P_n(\cos\theta_b)}\right]^2\right\} \qquad (10)$$

Note that if $\sigma = \infty$, we have

$$p(x) = p_e(x) + \frac{B_o^2}{8\pi x^4} = p_e(x) + \frac{[B_e(x)]^2}{8\pi}$$

where $B_e(x)$ is the external magnetic field strength. This is the condition one would expect were there no magnetic field inside the sheet. We now see that the departure from this condition is entirely due to finite resistivity effects since these effects introduce an internal magnetic field.

Evaluating equation (10) at $x = 1$, we can solve for n as a function of the pressure difference across the sheet.

$$n = \sqrt{R_{em}^2 u_1^2 (1-\alpha) - R_{em} \frac{d}{dx}(x^2 u)\Big|_{x=1}} \qquad (11)$$

where $u_1 = u(x=1)$ and $\alpha = (E-1)/\phi \cdot E = P_s(1)/P_e(1)$, with $P_s(1)$ being the pressure in the sheet at $x = 1$ and $P_e(1)$ the pressure just outside, and $\phi = B_o^2/8\pi P_e(1)$ represents the ratio of magnetic to gas pressure outside the sheet at the reference level. A cursory examination of equation (11) might suggest that $\delta \propto 1/R_{em}$ (since $\delta \propto 1/n$), yielding an extremely small value for the sheet thickness. However, the velocity terms u_1 and $d/dx(x^2u)|_x = 1$ are expected to be very small. Also, if $\sigma = \infty$, $P_s(1) = P_e(1) + B_o^2/8\pi$ [*Pneuman and Kopp*, 1971] so that $\alpha = 1$. In our case α will not be exactly one but extremely close to one since $R_{em} >> 1$. Hence, the right-hand side of equation (11) will be significantly reduced by these considerations, resulting in an increased sheet thickness.

To evaluate α, u_1, $d/dx(x^2u)|_x = 1$, it is necessary to calculate the flow along the sheet. To do this, we must take into account the radial **j**X**B** force, which because of the high current density is *not* small. (It will be shown that the radial **j**X**B** force is actually independent of the electrical conductivity close to the sun.) The radial momentum equation, once the density is eliminated via the continuity equation, then becomes approximately

$$\left(u - \frac{1}{u}\right)\frac{du}{dx} = \frac{2}{x} - \frac{\psi}{x^2} - \frac{2n^2\phi}{(1+\alpha\phi)u_1 R_{em} x^4} \qquad (12)$$

where $\psi = GM_\odot / r_o V_s^2$. In deriving equation (12), we have assumed that

$$\exp\left[-\frac{n(n+1)}{R_{em}} \int_1^x \frac{dx}{x^2 u}\right] \approx 1$$

Equation (12) without the last term is the solar wind equation used by *Parker* [1958] for an isothermal corona. The critical point in that case is seen to be at $x_c = \psi/2$ where $u = 1$. The last term here is the radial component of the magnetic force, which is negative and hence acts to retard the expansion. As a result, the critical point is moved outward and occurs where the left-hand side of equation (12) vanishes, yielding

$$x_c^2\left(x_c - \frac{\psi}{2}\right) = \frac{n^2\phi}{R_{em} u_1 (1+\alpha\phi)} \qquad (13)$$

For no field, $\phi = 0$ and equation (13) reduces to $x_c = \psi/2$ [*Parker*, 1958].

We now integrate equation (12) to obtain

$$\frac{u^2}{2} - \ln u - 2\ln x - \frac{\psi}{x} - \frac{2n^2\phi}{3R_{em} u_1 (1+\alpha\phi) x^3} = const$$

which, evaluated both at the critical point and at $x = 1$, yields

$$\frac{1}{2} - 2\ln x_c - \frac{\psi}{x_c} - \frac{2n^2\phi}{3u_1 R_{em}(1+\alpha\phi)x_c^3}$$

$$= \frac{u_1^2}{2} - \ln u_1 - \psi - \frac{2n^2\phi}{3u_1 R_{em}(1+\alpha\phi)} \qquad (14)$$

Equations (13) and (14) represent two coupled equations for x_c and u_1. However, the quantities n and α are also functions of u_1 and must be evaluated.

To do this, consider the portion of the streamer below r_o which leads up into the sheet. Here the flow is along the field and no magnetic force exists along field lines.

Conservation of momentum along the field then shows that

$$P_s(1) = p_o \exp\left[-\left(1 - \frac{R_\odot}{r_o}\right) + \frac{1}{2}\left(u_0^2 - u_1^2\right)\right] \quad (15)$$

where p_o is the pressure at the solar surface ($r = R_\odot$) and u_o is the velocity there. Noting that $u_o \ll u_1$, we can write equation (15) in the form

$$p_s(1) = p_c^* \exp\left(-\frac{u_1^2}{2}\right)$$

where $p_c{}^*$ is the pressure at the neutral point that would be obtained for $\sigma = \infty$ (the hydrostatic value). Now we have

$$\alpha = \frac{E-1}{\phi} = \frac{[P_s(1)/P_e(1)]-1}{\phi}$$

$$= \frac{[p_c^*/p_e(1)]\exp(-u_1^2/2) - 1}{\phi}$$

But $p_c{}^*/p_e(1) = 1 + \phi$ [*Pneuman and Kopp,* 1971]. Hence

$$\alpha = \frac{(1+\phi)\exp(-u_1^2/2) - 1}{\phi}$$

and, since $u_1 \ll 1$, we can expand the exponential in the numerator and obtain

$$\alpha \approx 1 - \left(\frac{1+\phi}{2\phi}\right) u_1^2$$

Using equation (12), we also can evaluate $d/dx(x^2u)|_{x=1}$ and find

$$\frac{d}{dx}(x^2u)\Big|_{x=1} = \psi u_1 + \frac{2n^2\phi}{R_{em}(1+\phi)}$$

Equation (11) now reduces to

$$n = \sqrt{R_{em}\left(\frac{1+\phi}{1+3\phi}\right) u_1 \left[R_{em} u_1^3 \left(\frac{1+\phi}{\phi}\right) - \psi\right]} \quad (16)$$

and, after some manipulation, equations (13) and (14) become

$$\left(\frac{1+\phi}{2\phi}\right) R_{em} u_1^3 = \psi + \left(\frac{1+3\phi}{\phi}\right) x_c^2 \left(x_c - \frac{\psi}{2}\right) \quad (17)$$

$$\frac{u_1^2}{2} - \ln u_1 = \frac{2}{3} x_c^2 \left(x_c - \frac{\psi}{2}\right) - 2\ln x_c + \psi - \frac{1}{6} \quad (18)$$

Noting that $u_1^2/2 \ll |\ln u_1|$ we finally obtain a single transcendental equation for x_c.

$$\ln\left[\frac{2\psi_\phi + (1+3\phi)x_c^2(2x_c - \psi)}{R_{em}(1+\phi)x_c^6}\right] + x_c^2(2x_c - \psi) - \frac{2\psi}{x_c} + 3\psi - \frac{1}{2} = 0 \quad (19)$$

The numerical procedure is to solve equation (19) for x_c; then u_1 can be found from equation (17) or (18). The sheet thickness $\delta(\sim\pi r/4n)$ can then be evaluated from equation (16). Figure 2 shows a plot of δ at the coronal base as a function of temperature for various values of ϕ. For reasons to be discussed later, these values should properly be considered lower limits to the actual sheet dimensions.

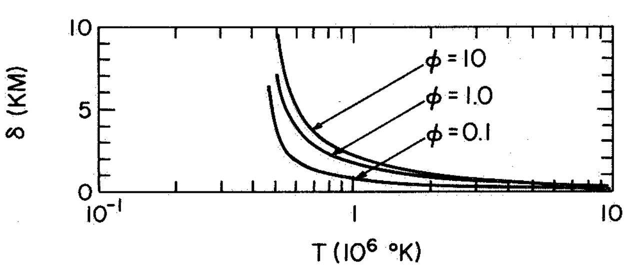

Figure 2. *Sheet thickness at the coronal base as a function of coronal temperature. These values are based upon an electrical conductivity unaffected by the magnetic field and should be considered lower limits.*

We now summarize the solutions (approximate) for the pertinent properties inside the sheet using the approximations $R_{em} \gg 1$ and

$$\exp\left[-\frac{n(n+1)}{R_{em}} \int_1^x \frac{dx}{x^2u}\right] \approx 1$$

$$B_r = \frac{B_o}{x^2} \sin\left[\left(n + \frac{1}{2}\right) \theta + \frac{\pi}{4}\right] \qquad (20)$$

$$B_\theta = \frac{-B_o n}{R_{em} x^3 u} \cos\left[\left(n + \frac{1}{2}\right) \theta = \frac{\pi}{4}\right] \qquad (21)$$

$$J = \frac{-B_o n}{4\pi r_o x^3} \cos\left[\left(n + \frac{1}{2}\right) \theta + \frac{\pi}{4}\right] \qquad (22)$$

$$p = p_e(x) + \frac{B_o^2}{8\pi x^4} \cos^2\left[\left(n + \frac{1}{2}\right) \theta + \frac{\pi}{4}\right] \qquad (23)$$

The velocity u is given by the equation

$$\frac{u^2}{2} - \ln u = \frac{u_1^{\ 2}}{2} - \ln u_1 + 2 \ln x - \psi \left(1 - \frac{1}{x}\right) - \frac{2n^2 \phi}{3u_1 R_{em}(1 + \phi)} \left(1 - \frac{1}{x^3}\right) \qquad (24)$$

Figure 3 shows the expansion velocity both inside and outside the sheet as a function of radial distance. The

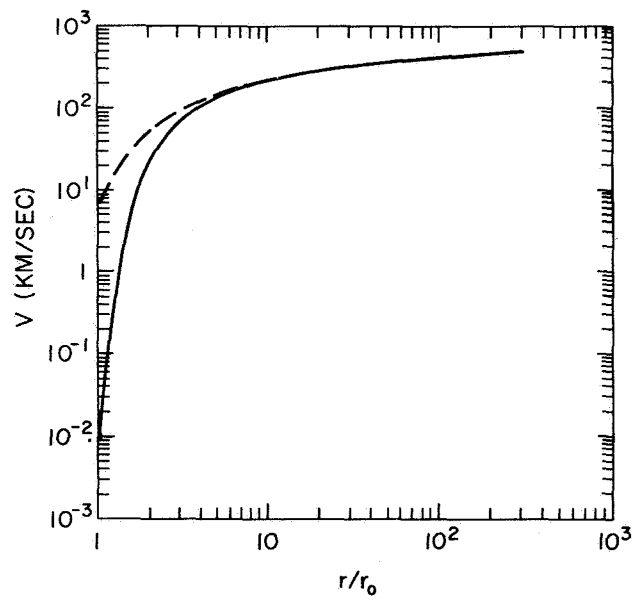

Figure 3. *Variation of expansion velocity with radial distance for $\phi = 1$, $r_o = 2 R_\odot$, and $T = 1.0 \times 10^6$ °K. The solid curve represents the velocity inside the sheet whereas the dashed curve is the velocity just outside. Note that the flow in the sheet near the base is extremely small due to the inhibiting effect of the magnetic field.*

small velocities in the sheet near the coronal base reflect the retarding force of the internal magnetic field. The density enhancement for various field strengths is shown in figure 4. As expected, the enhancement close to the sun increases with increasing field strength (increasing ϕ) In this model, the enhancement falls to zero at large distances; however, this behavior may be due to the omission of thermal conduction effects in the model. It is expected that thermal conduction will be inhibited in the sheet by the transverse magnetic field, whereas conduction is essentially unimpeded outside. As shown by *Pneuman and Kopp* [1970], the inclusion of this effect could tend to maintain significant density enhancements at large distances.

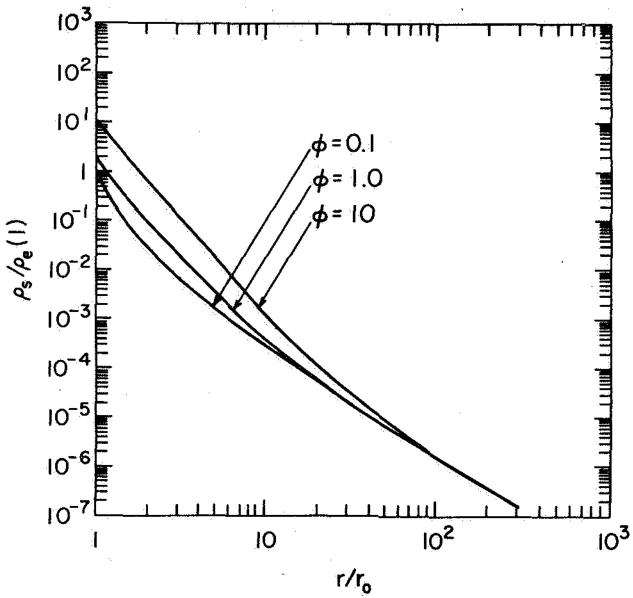

Figure 4. *Density enhancement (ratio of density in the sheet to that outside) as a function of radial distance for the cases $\phi = 0.1$, 1, and 10.*

It is of special interest to note the dependence of the various physical quantities on the electrical conductivity R_{em}. Solutions of equation (19) over a wide range of magnetic Reynolds numbers reveals that the location of the critical point (x_c) is not a very sensitive function of R_{em}. For example, for coronal temperatures in excess of 10^6 °K, x_c only varies from 1 to about 4 for $1 < R_{em} < 10^{14}$. Equation (17) then shows that $u_1 \propto R_{em}^{-1/3}$, and from equation (16) we have $n \propto R_{em}^{1/3}$. This means that the sheet thickness δ varies as

$$\delta \propto R_{em}^{-1/3}$$

We can now use equations (21) and (22) to evaluate the transverse field in the sheet (B_θ) and the electric current density J. Close to the sun, $u \approx u_1 \propto R^{-1/3}$. Hence, near $x = 1$, we have (since $n \propto R_{em}^{1/3}$).

$$B_\theta \propto R_{em}^{-1/3} \qquad x \approx 1$$

At larger distances, u is of order 1 and

$$B_\theta \propto \frac{R_{em}^{-2/3}}{x^3} \qquad x >> 1$$

The current J has the dependence

$$J \propto \frac{R_{em}^{1/3}}{x^3}$$

Consequently, the radial **J**X**B** force ($\propto JB_\theta$) is independent if R_{em} close to the sun but varies as $R_{em}^{-1/3}$ at large x, signifying that the magnetic stresses are important at lower levels but do not significantly affect the expansion at large distances.

DISCUSSION

Since the neutral sheet model presented in the previous sections is based upon the MHD approximation, it is appropriate now to examine its validity in structures whose lateral dimensions may be of the order of a mean free path. For example, a coronal temperature of 10^6 °K and an electron density of $10^8/cm^3$ corresponds to a mean free path of 100 km, which is larger than the dimensions shown in figure 2. However, the Larmor radius, probably the more appropriate dimension in this case, is only 20 cm for a field strength of 1 gauss. This dimension is small relative to the thickness so that electrons are very effectively tied to the field lines. For this reason, ion-electron collisions will be frequent making the MHD approximation a reasonable initial approach to the problem.

The dimensions shown in figure 2 should also properly be considered as lower limits since they are based on an electrical conductivity unaffected by the magnetic field. From equation (21) the transverse field B_θ at the center of the sheet at $x = 1$ for $B_O = 1$ gauss, $r_O = 2R_\odot$, and a temperature of 10^6 °K ($R_{em} = 4.5 \times 10^{14}$) is found to be only 3×10^{-5} gauss and decreases outward rapidly. The product of the cyclotron frequency and collision time for electrons in this field is about 70. Hence, R_{em} could be reduced by a factor as great as 2×10^{-4}. This consideration would increase the sheet dimension by the factor 17 yielding widths in the inner corona in the order of 100 km. This figure corresponds to a thickness of 10,000 km at 1 AU. Another important consideration may be various forms of microturbulence and instabilities in the plasma, which would tend to increase the "effective" collision frequency, resulting in a further reduction in electrical conductivity.

With regard to the generating mechanism for type III radio emission, we should note that the neutral sheet described here contains a weak but finite transverse magnetic field. This field would have a retarding effect on electrons traveling outward along the sheet. One might well ask whether the concept of electrons moving outward at velocities of the order $0.3c$ across a magnetic field of 10^{-5} gauss is realistic. On the other hand, stabilization of the beam in such a configuration seems more possible here than in the ambient corona since the rapid increase in field strength outward from the center of the sheet should have a confining effect on the electrons.

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DISCUSSION

R. A. Kopp It's always nice to not worry about the convergence problem by iterating just once, but it seems like your thickness is varying by three orders of magnitude from the assumption of radial. Are you sure this is conclusive? Should't you work it out?

G. W. Pneuman I would say it is not a conclusive result, no.

F. W. Perkins As far as I can see, the reason why we require finite conductivity is because you assume that there was flow out. I suspect there is a possibility of having a static solution, too, with infinite conductivity. Is the flow really observed or necessary?

G. W. Pneuman I believe it is necessary to satisfy the condition of almost zero pressure at infinity.

T. G. Cowling I would like to ask, in view of the background to these latest contributions, whether in fact we are approaching the idea that the layer that you get at the boundary between two sectors is to be regarded as a shear layer, the equivalent of the earthquake boundary we've heard so much about? Is it a thing where the field virtually becomes detached, field from the one side is to be regarded as detached from the field on the other side? Is that a picture we have come to? And is there a genuine reason for that?

G. W. Pneuman I'm not sure. In this particular model I would consider it that way. I would like to look at it as the loops, the low down loops, actually being drawn up inside the sheet so that actually there are never really any open field lines inside the sheet; sort of an extension of the lower part of the streamer.

Unidentified You'd still have a discontinuity involved above and below?

G. W. Pneuman No, there is no discontinuity.