HEATING SOLAR CORONAL HOLES

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ABSTRACT

It has been shown that the coronal hole, and the associated high-speed stream in the solar wind, are powered by a heat input of the order of $5 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$, with most of the heat injected in the first $1-2 R_\odot$, and perhaps $1 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ introduced at distances of several solar radii to provide the high speed of the issuing solar wind. The traditional view has been that this energy is obtained from Alfvén waves generated in the subphotospheric convection, which dissipate as they propagate outward, converting the wave energy into heat. This paper reviews the generation of waves and the known wave dissipation mechanisms, to show that the necessary Alfvén waves are not produced under the conditions presently understood to exist in the Sun, nor would such waves dissipate significantly in the first $1-2 R_\odot$ if they existed. Wave dissipation occurs only over distances of the order of $5 R_\odot$ or more.

The alternative to wave heating is the activity of the small-scale magnetic fields—the network and intranetwork fields. Existing estimates indicate an adequate energy input at the base of the coronal hole and a sufficient generation of Alfvén waves to account for the distant heat input. It appears, then, that the network activity is the principal energy source for the coronal hole and hence is the basis for creating the heliosphere.

Subject headings: radiative transfer — Sun: corona — Sun: solar wind — wave motions

1. INTRODUCTION

This paper addresses the theoretical problem posed by the heat input to the coronal hole, producing the high-speed streams (400–600 km s$^{-1}$) in the solar wind. Withbroe & Noyes (1977) estimate the heat input at $5-8 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ at the Sun, with most of the energy going into expansion of the coronal gas ($dQ \approx dV$) to produce the solar wind (see reviews by Hollweg 1978; Holzer & Leer 1980; Pneuman 1986; Mullan & Waldron 1987; and the specific cases treated by Hollweg 1973; Hammer 1982a, b; Jatenco-Pereira & Opher 1989; and Wendel 1989). The recent analysis of the observed structure of the coronal hole by Withbroe (1988) shows that a total heat input of $5 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ with the major portion deposited in the first one or two solar radii above the surface of the Sun ($1 < r < 3 R_\odot$). The final asymptotic high-speed of the expansion requires approximately $1 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ of the total to be deposited over extended radial distances beyond the sonic point in the wind (Parker 1964a, b, 1965; Barnes 1969; Leer & Holzer 1979; Holzer & Leer 1980; Leer, Holzer, & Flá 1982). So if the total input is $5 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$, some $4 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ is deposited in the near coronal hole and $1 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ in the far coronal hole.

The traditional idea has been that the heating is a result of the dissipation of waves propagating outward from their place of origin in the photospheric convection. We note, then, that Athay & White (1979a, b) and Leer et al. (1982) infer an upper limit on the outward Alfvén wave flux in the coronal hole of about $1 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ from their analysis of the observations. This upper limit allows only enough energy to supply the distant corona. Note, then, that Habbal & Leer (1982) and Flá et al. (1984) have shown that Alfvén and fast-mode waves are damped by electron thermal conduction over distances of the order of $10 R_\odot$ (see also Barnes 1969). So if the necessary Alfvén wave flux of about $1 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ is present, it provides adequate heat to the distant coronal hole.

Consider, therefore, the properties of such waves. A magnetic field $B = 10$ G in a coronal hole where the temperature $T = 1.5 \times 10^6$ K and the number density $N = 10^8$ atoms cm$^{-3}$ (each atom of mass $M = M_p$) leads to an Alfvén speed $C = 2 \times 10^8$ cm s$^{-1}$ in association with a sound speed $c = 2 \times 10^7$ cm s$^{-1}$. The agitation observed in the photosphere suggests characteristic wave periods $\tau = 50–300$ s. A Kolmogorov spectrum in the photosphere produces agitation at shorter periods, but with the kinetic energy density declining in direct proportion to $\tau$.

The magnetic field dominates the gas in the coronal hole, with $B^2/8\pi \approx 10^7$ p $\approx 4$ dyn cm$^{-2}$, where $p$ is the gas pressure $2 N k T$ (for ionized hydrogen). The slow-mode magnetohydrodynamic wave is essentially a sound wave with the displacement of the gas constrained to the direction of the magnetic field. The fast-mode wave propagating near the direction of the field is a compressible Alfvén wave, with $|\delta p + B^2/4\pi| \ll |\delta p|$, where $\delta p$ and $\delta B$ represent the variation of $p$ and $B$ through the wave.

A fast-mode wave with a period $\tau = 10^2$ s has a wavelength $\lambda = 2 \times 10^1$ cm and a wave number $k = 2\pi/\lambda = 3 \times 10^{-10}$ cm$^{-1}$. The electron thermal conductivity $\kappa$ is equal to $10^{-8} T^{5/2}$ ergs cm$^{-1}$ s$^{-1}$ K in a collision-dominated plasma. The thermometric conductivity $K$ (cm$^2$ s$^{-1}$) is $(\gamma - 1)k/\gamma N k = 0.4k/\gamma N k$ for $\gamma = 5/3$ (in which $k$ represents the Boltzmann constant). The characteristic damping times of the thermal energy associated with the wave number $k$ is $t_D = \frac{1}{2} k^2/k$, which yields $t_D = 4$ s. However, the thermal energy in the fast-mode wave constitutes only a fraction $O(c^2/C^2) = 10^{-2}$ of the total energy, so the effective damping time for the wave is about $10^2$ times longer, of the order of $4 \times 10^2$ s. With an Alfvén speed of $2 \times 10^8$ cm s$^{-1}$, this estimate of the damping length is of the order of $8 \times 10^9$ cm $\approx 1 R_\odot$. The more detailed study by Habbal & Leer (1982), recognizing the collisionless nature of the outer corona, suggests damping over somewhat longer characteristic distances, of the order of $5 R_\odot$.
(see also the extended discussions by Holzer, Flà, & Leer 1983; and Barnes 1969, 1979).

It seems clear, then, that fast-mode waves and Alfvén waves are damped, but only over radial distances of $5-10 R_\odot$. So if a wave flux of the order of $1 \times 10^4$ ergs cm$^{-2}$ s$^{-1}$ with periods of the order of $10^2$ s exists in the corona, it should provide the energy at large distance that is necessary to boost the thermal expansion velocity to 500–600 km s$^{-1}$.

As a point of information, a wave with rms velocity $\langle v^2 \rangle^{1/2}$ and a group velocity $V$ carries an energy flux $I$ given by

$$I = \rho \langle v^2 \rangle V \text{ ergs cm}^{-2} \text{ s}^{-1}. \quad (1)$$

Note that for a transverse wave the vibrations $\langle v^2 \rangle^{1/2}$ may be in each of the two directions perpendicular to the direction of propagation, so that in any one transverse direction the rms velocity is $\langle v^2 \rangle^{1/2}$. In the present instance, $V = C \cong 2 \times 10^8$ cm s$^{-1}$ in the coronal hole, and $I \cong 1 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$, so that $\langle v^2 \rangle^{1/2} \cong 12$ km s$^{-1}$. In the chromosphere, where $T = 7 \times 10^4$ K and $N = 2 \times 10^{10}$ atoms cm$^{-3}$, the Alfvén speed $C \cong 1.2 \times 10^3$ cm s$^{-1}$, and the sound speed $c = 1.1 \times 10^5$ cm s$^{-1}$, so that the same $I$ and $B$ yield $\langle v^2 \rangle^{1/2} \cong 3.5$ km s$^{-1}$ for Alfvén waves. In the photosphere, where $T = 5.6 \times 10^4$ K and $N = 10^{17}$ atoms cm$^{-3}$, the Alfvén speed $C = 6 \times 10^5$ cm s$^{-1}$, and the sound speed $c = 0.9 \times 10^6$ cm s$^{-1}$, so that $\langle v^2 \rangle^{1/2} \cong 7 \times 10^2$ km s$^{-1}$. Note, then, that in the photosphere the rms transverse motions are as large as the Alfvén speed. This is a wave of large amplitude $A$. For a wave number $k$, we have $kA = \langle v^2 \rangle^{1/2} C \cong 2 \times 10^3$.

The amplitude scales as $I^{1/2}$, of course, so that $I = 5 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ to supply the entire energy, would require $kA \cong 3$. And if the high reflectivity (~0.7) (Jeffrey 1966; Hollweg 1984, 1985) at the chromosphere-corona transition is taken into account, across which the Alfvén speed and the sound speed both jump by more than a factor of 10, the necessary $kA$ in the photosphere becomes at least five. The wave forms would be hairpin loops, like the meandering of a river over a broad floodplain.

Now the problem is twofold. First of all, it is doubtful that the photospheric convection produces an Alfvén wave flux providing $1 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ in the corona, necessary to supply the outer coronal hole, particularly when the high reflectivity of the chromosphere-corona transition region is included. Second, even if, somehow, there was an Alfvén wave flux as large as $5 \times 10^4$ ergs cm$^{-2}$ s$^{-1}$ in the coronal hole, there is no known way by which such waves with period $T$ of 50 s or more can damp in the available length of $1-2 R_\odot$ to provide the necessary heat input to the near coronal hole.

These two points are developed in the succeeding sections of this paper. In particular, we examine every dissipation scheme of which we are aware to show that the dissipation of waves beyond about $5 R_\odot$ is to be expected, but with little effect in the first $1-2 R_\odot$, whereas Withbroe’s (1988) analysis requires the opposite.

This forces us to the conclusion that the coronal hole is not heated principally by the dissipation of waves produced in the subphotospheric convection. The only available alternative is the activity of the network and intranetwork magnetic fields. The idea that these fields are the prime energy source for the coronal hole was first advanced by Martin (1984, 1988) and has been expounded more recently by Porter et al. (1987; Porter & Moore 1988). Our purpose here is to show that the idea is evidently correct, insofar as it can be checked and substantiated by observations, and to show that there is no viable alternative within the framework of our present (limited) knowledge of the physics of the Sun. It follows that observational studies of the behavior of the network and intranetwork magnetic fields take on a central role in the science of the coronal hole. The complex activity of the small-scale magnetic fields has already been established (Martin 1984, 1988; Porter et al. 1987; Wang, Zirin, & Ai 1990), with rapid changes in field patterns, the appearance and disappearance of flux elements, and vigorous microflaring. The subject is of interest in its own right for what it can tell us about the magnetic activity above and below the visible surface. However, the scientific study of the network and intranetwork activity takes on a new dimension because, as the principal heat source to the solar corona hole, it becomes the agent primarily responsible for creating the heliosphere (see discussion in Parker 1987b). It is curious that the smallest scale activity appears to be responsible for the largest scale effect.

The remainder of the paper develops the foregoing arguments in detail. The next section looks into the production and transmission of waves from the subphotospheric convection. The section after that (§ 3) is a meticulous (and necessarily tedious) examination of the many known wave-damping mechanisms, to establish the inapplicability of the traditional view of coronal heating. The following section (§ 4) explores the alternatives.

## 2. Waves in the coronal hole

Theoretical studies of the propagation of plane acoustic waves, gravitational acoustic waves, and fast- and slow-mode magnetohydrodynamic waves (Osterbrock 1961; Stein 1968; Stein & Schwartz 1972, 1973; Stein & Leibacher 1974; Schwartz & Stein 1975) indicate that these waves are either dissipated by shock formation or refracted away from the vertical before reaching the corona. Further, Athay & White (1978, 1979a; see also Sturrock & Uchida 1981) infer an upper limit of $1 \times 10^4$ ergs cm$^{-2}$ s$^{-1}$ on acoustic waves from observations of the chromosphere. At most, then, it is only Alfvén waves that reach the corona (Kuperus, Ionson & Spicer 1981; Priest 1982), and even the Alfvén waves are attenuated by the abrupt increase in the Alfvén speed by more than a factor of 10 (over about $10^8$ cm) in passing from the chromosphere to the corona (Hollweg 1984, 1985). The reflection coefficient is of the order of 0.7 for wave periods of $10^2$ s. An et al. (1989) have shown that for wave periods of $10^3$ s or more, the reflection becomes so strong as to provide trapping and resonance$^1$ (see also Moore & Musielak 1991).

There are resonance effects that enhance the transmission of Alfvén waves into the corona (Hollweg 1981b, 1984; Zugsda & Locans 1982), but the effects arise only in the bipolar magnetic fields of (small or large) active regions or in spicules (Sterling & Hollweg 1984), so they are not applicable to wave transmission in the unipolar field of a coronal hole.

Hollweg (1982a) examines the coronal heating by shock-on shocks produced by Alfvén waves with long periods of 1–2 hr. However, the power expected in Alfvén waves of such long period is not adequate for the requirements of the coronal hole. Flà et al. (1984) have pointed out the possibility that the agitation in the large bipolar magnetic fields of active regions

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$^1$ There is a comparable upper limit $4\pi\eta/\rho g$ of about 600 s on the period of sound waves for vertical propagation in an isothermal chromosphere, where the sound speed $c \cong 2 \times 10^6$ cm s$^{-1}$ and $g = 2.7 \times 10^4$ cm s$^{-2}$ with $\gamma = 5/3$ (cf. Musielak 1990).
may introduce substantial fluxes of fast-mode waves into the neighboring coronal holes. However, the coronal holes and fast wind streams appear to exist independently of the active regions (which essentially vanish at sunspot minimum), so the effect cannot be the staple energy supply we are seeking.

It would appear, then, that Alfvén waves with periods of 50–300 s, characteristic of the granule agitation (on a scale $l \approx 500$ km and a characteristic velocity $v \approx 10^3$ cm s$^{-1}$), are the only candidates for photospheric wave heating in the coronal hole. All other modes are relatively ineffective either in their creation or in their propagation into the corona.

The theory of the generation of magnetohydrodynamic waves in the subsolar photosphere has taken substantive form in recent years (Mursel & Rosner 1987, 1988; Mursel, Rosner, & Ulmschneider 1989; Collins 1989a, b, 1991). As already noted, in the photosphere the sound velocity is $10^3$ cm s$^{-1}$ and the Alfvén speed $C$ in 10 G is $6 \times 10^3$ cm s$^{-1}$ (for $N_M = 2 \times 10^{10}$ g cm$^{-3}$). The fast-mode wave is essentially a sound wave and the slow-mode wave is an oblique Alfvén wave. An Alfvén wave carrying $1 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ has an rms velocity $\sqrt{v^2/2} = 7 \times 10^2$ cm s$^{-1}$ in each direction perpendicular to the magnetic field. As already noted, this is approximately equal to the Alfvén speed, so that the waves are of large amplitude.

Collins (1991) points out that the generation of Alfvén waves is relatively weak in the photospheric convection, because the Alfvén speed $C \approx 6 \times 10^3$ cm s$^{-1}$ is small compared to the characteristic phase speed ($\sim 10^4$ cm s$^{-1}$) of the principal Fourier components of the granule convection. To state this differently, an Alfvén wave with a period $\tau$ as long as the characteristic granule period $\tau = 5 \times 10^2$ s has a wavelength of only $3 \times 10^8$ cm, equal to 0.06 $R_\odot$. Shorter wave periods of 50–300 s give corresponding shorter wavelengths, with 50 s providing a wavelength of 0.006 $R_\odot$. These are small compared to $l$.

It was pointed out in §1 that the wave flux of $5 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ to heat the outer coronal hole requires Alfvén waves of large amplitude $A$ in the photosphere, with $kA \approx 5$. But turbulent eddies have a characteristic life $\tau \approx t_e$, with a characteristic displacement amplitude $A = vt_e$ and wave number $k = O(l)/l$ so that $kA = O(1)$. That is to say, the individual eddy provides deformation of a field line that is only about as deep as it is broad. It is not expected to supply the deep hairpin loops required to provide a wave flux of $5 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ in the corona. Collins’s formal calculations bear out the weak generation of Alfvén waves. If the Alfvén speed were as large as the convective velocity, it would be a different story, of course, because there would then be a good impedance match to the waves.

It is interesting to note that the fibril structure of the photospheric magnetic field may play an important role here, providing an Alfvén speed of $6 \times 10^3$ cm s$^{-1}$ in the fibril field of $10^3$ G. The fibril admits surface waves with phase velocities ranging downward from about $4$ km s$^{-1}$ (cf. Parker 1974; Roberts 1976b), thereby providing a better match to the granule velocities. What is more, the wave modes are convectively overstable and may therefore be a direct source of upward propagating Alfvén waves. (Parker 1974, 1975; Roberts 1976a, b, 1983). It is difficult on theoretical grounds to state just how strong such fibril waves might be. Now the Alfvén transit time from the photosphere to the corona is of the order of $10^2$ s, depending on the particular model used for the magnetic fibril (cf. Simon, Weiss, & Nye 1983; Fiedler & Cally 1990; Parker 1991). The essential point is that whatever model one uses, the transit time is less than the characteristic granule life $l/v \approx 5 \times 10^2$ s. Hence the fibril moves in a quasi-rigid form as it is buffeted by the photospheric convection, providing transverse motions at the corona comparable to the transverse convective motions of $v = 1$ km s$^{-1}$ in the photosphere. If it is argued that the fibril is caused to lash back and forth, as a consequence of eddies of scale $l$ and velocity $v$ at the photosphere, then the semirigid fibril might sweep back and forth with a velocity $\nu z/l$ at a height $z$ above the photosphere. With $l = 500$ km and $v = 10^2$ cm s$^{-1}$, this yields transverse motions of $10$ km s$^{-1}$ at a height $z = 5 \times 10^4$ cm in the coronal hole. But $10$ km s$^{-1}$ just barely supplies the $1 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ to heat the distant coronal hole. It falls far short of the total $5 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$.

3. WAVE DISSIPATION

Suppose that there were Alfvén waves (of unknown origin) with characteristic periods of 50–300 s propagating outward in the coronal hole. Consider whether such Alfvén waves might be dissipated in the first one or two solar radii, requiring some mechanism that is a factor of 5 or 10 more effective than electron thermal conduction. Phase mixing has been a popular idea for dissipating Alfvén waves (Tataronis & Grossman 1973; Sakurai & Granik 1984; Nocera, Leroy, & Priest 1984; Mok & Einaudi 1985; Steinolfson 1985; Lee & Roberts 1986). But phase mixing is expected to dissipate Alfvén waves only over distances of $10 R_\odot$ or more (Nocera et al. 1984), and, in any case, phase mixing requires an ignorable transverse coordinate, which occurs only if the medium is uniform in that coordinate. In view of the extreme inhomogeneity of the transition region (Kopp & Kuperus 1968; Orrall et al. 1990), a sufficiently uniform corona is dubious. Hence we cannot expect phase mixing to be effective (Parker 1991). However, by the same token, we expect to find the resonance of a wave with frequency $\omega$ and wave number $k(z)$ at surfaces extending along the field $B(x, y)$ where the local Alfvén speed $C(x, y, z)$ is equal to $\omega/k(z)$. This provides interest in Alfvén internal wave structures including longitudinal oscillations that introduce dissipation (Grossmann & Tataronis 1973; Chen & Hasegawa 1974; Tataronis 1975; Kappraff & Tataronis 1977; Wentzel 1979a, b; Sakurai & Granik 1984; Hollweg 1984, 1985, 1987a, b; Donnelly, Clancy, & Cramer 1985; Lee & Roberts 1986; Davila 1987; Einaudi & Mok 1985, 1987; Hollweg 1987a, b; Grossman & Smith 1988; Strauss & Lawson 1989; An et al. 1989; Mok & Einaudi 1990; Hollweg et al. 1990, and references therein). Lee & Roberts describe the mode conversion involved in the resonance in some detail, noting that the energy of the fast-mode wave is converted into local longitudinal oscillations at the resonance surface, and that the longitudinal oscillations are eventually damped by viscosity so that the energy is converted to heat in distances of 5–20 $R_\odot$. The final dissipation rate is comparable in magnitude to the effect of electron thermal conduction.

There are other ways in which Alfvén or fast-mode waves may be damped in a low $\beta$ plasma, of course. Similon & Sudan (1989) have shown how the expected random walk of the field lines in a static field (Jokipii & Parker 1968, 1969a, b) disperses Alfvén waves propagating along the magnetic field of a bipolar active region. Unfortunately the stochastic field lines in a
coronal hole do not accomplish the same dispersion because the random transverse components of the field that produce the random walk of the field lines are themselves Alfvén waves propagating outward along the field. The stochastic component is part of the wave field rather than a standing structure that scatters the wave field.

The formation of tangential discontinuities, as a consequence of the quasi-static deformation of the bipolar fields of active regions, appears to be the source of dissipation of magnetic field responsible for the active X-ray corona (Parker 1981b, 1983b, 1987c, 1988). But this, too, is inapplicable to the coronal hole, because, again, the transverse excursions of the field lines responsible for causing the discontinuities propagate away from the Sun as Alfvén waves, taking any potential discontinuities with them.

It must be kept in mind that the amplitude of the Alfvén waves in the corona is small. It was pointed out in §1 that, while the amplitude in the photosphere is large, in the corona the relative amplitude, measured by \( \langle B^2 \rangle^{1/2} / B \) and \( \langle v^2 \rangle^{1/2} / C \), is only \( O(10^{-2}) \). It follows that the nonlinear steepening of wavefronts and mode-mode coupling are weak and can be neglected in our search for rapid dissipation of the waves in the first ten wavelengths (1–2 \( R_\odot \)).

Zhu (1991) has pointed out that the interaction of Alfvén waves with the curvature of the magnetic field, producing slow-mode waves, is a first-order effect. The effect is strong when \( \beta = O(1) \) and the wavelength is comparable to the radius of curvature of the field, but too slow when \( \beta \) is only of the order of \( 10^{-2} \) in the coronal hole.

One sometimes hears the suggestion that plasma turbulence causes anomalous resistivity, which rapidly damps Alfvén waves, etc. Unfortunately, strong plasma turbulence arises only when the electron conduction velocity \( u \) becomes comparable to some significant fraction of the ion thermal velocity \( w \) of about \( 2 \times 10^6 \) cm s\(^{-1}\) (Buneman 1958, 1959; Stringer 1964; Kadomtsev 1965; Sagdeev 1967; Kalinin et al. 1970). The electron drift instability (Haerendel 1977) appears with the smallest electron drift velocity, but the stronger turbulence arises from the lower-hybrid drift instability, when \( u \) exceeds \( (m/M)^{1/4} w \approx 0.15w \) (Coroniti & Evitt 1977; Huba, Glad, & Papadopoulos 1977; Tanaka & Sato 1981). Ion acoustic turbulence appears when \( u \) exceeds \( w \) (Hamburger & Friedman 1968; Friedman & Hamburger 1969; Diamond et al. 1984; Haerendel 1990). Hamburger & Friedman find that the effective electron collision frequency is approximately one-tenth of the electron plasma frequency \( v_p \) (or \( 10^{-2} \omega_p \), where \( 2\pi v_e = \omega_p = [4\pi Ne^2/m]^{1/2} \)). If the drift velocity \( u \) exceeds the electron thermal velocity \( w(m/M)^{1/2} \), then the effective electron collision frequency increases until it is comparable to the plasma frequency \( v_p \) (or \( 10^{-1} \omega_p \)), because in this case \( u \) exceeds the Langmuir wave velocity, so that electrostatic plasma oscillations are excited. In fact, at this level, it becomes possible to establish electric double layers (Alfvén & Carlquist 1967; see Heyvaerts 1981, and references therein), with consequences for particle acceleration. This interesting situation is not achieved in any situation of which we are aware in the dense plasmas of the solar atmosphere. The scale of variation of the field would have to be small compared to the thermal ion cyclotron radius. The effect appears, however, in the very tenuous plasmas and thin current sheets in the terrestrial magnetosphere where it evidently plays an essential role in producing the aurorae (Temerin et al. 1982; Stenzel, Gekelman & Wild 1983). It appears that the collision frequency probably never exceeds \( 10^{-2} \omega_p \) in the magnetic fields at the Sun, except perhaps in solar flares (Heyvaerts 1981).

Now the electron conduction velocity \( u \) follows from Ampere's law, which gives the current density \( j = -Neu \) in terms of the curl of the magnetic field as

\[
Neu = c \Delta B / 4\pi \Delta l,
\]

where \( \Delta B = |\Delta B| \) is the change in \( B \) associated with \( \nabla \times B \) and an interval \( \Delta l \). It follows that only in extreme cases—quasi discontinuities—can we expect electron conduction velocities of sufficient magnitude to excite plasma turbulence. For suppose that \( \Delta B \) is as large as \( B = 10^6 \) G. Then we require \( \Delta l \approx 3 \times 10^4 \) cm in order that \( u \) be as large as the ion thermal velocity \( 2 \times 10^7 \) cm s\(^{-1}\). But the Alfvén waves that might heat a coronal hole would have an amplitude \( \delta B \) that is no more than \( 10^{-2} B_0 = 10^{-1} \) G, so that \( \Delta l \) must be as small as \( 3 \times 10^2 \) cm to admit the possibility of plasma turbulence. We know of no effect that would produce so small a transverse scale in \( B \) or in the wave field \( \delta B \). Indeed, the ion cyclotron radius is \( 2 \times 10^8 \) cm, so that \( \Delta l \approx 3 \times 10^2 \) cm is close to the extreme theoretical limit. Hence there is no reason to expect plasma turbulence and anomalous resistivity in coronal Alfvén waves.

Strauss (1988) has pointed out that the onset of the tearing-mode instability across a small-scale \( \Delta l \) produces an equivalent "hyperresistivity" with an associated diffusion coefficient

\[
Q = (\Delta B/B)^2 C/k^2 \Delta l \text{ cm}^2 \text{ s}^{-1},
\]

where \( k \) is the wave number. The characteristic damping time is

\[
t_D = 1/k^2 Q,
\]

\[
= (\Delta l/C)(B/\delta B)^3,
\]

so that the damping length is

\[
Ct_D = \Delta l(B/\delta B)^3.
\]

With \( \delta B = 10^{-2} B \), appropriate for waves in the solar corona, the damping length is of the order of \( 10^5 \) AU. A \( \Delta l \) as small as \( 10^5 \) cm yields a damping length of about \( 10^{13} \) cm, or about 1 AU. Stronger waves would be damped much more quickly, of course. But the effect does not appear to be effective in the waves of small amplitude expected in the coronal hole.

For the record, note that the ordinary conventional resistive diffusion coefficient \( \eta = c^2/4\pi Ne^2 \) cm\(^2\) s\(^{-1}\), valid for the expected small current densities, has a value of \( 3 \times 10^7 \) cm\(^2\) s\(^{-1}\) at \( 1.5 \times 10^7 K \) (\( a = 2 \times 10^7 T^{3/2} \text{ s}^{-1} \)). The characteristic resistive dissipation time across a scale \( l = 10^7 \text{ cm} \) is of the order of \( l^2/\eta l \), which becomes \( 10^9 \text{ s} \) \( \approx 30 \text{ yr} \).

The kinematic viscosity for ionized hydrogen has a value \( \nu = 1.2 \times 10^{-16} T^{3/2} / N \) cm\(^2\) s\(^{-1}\), equal to \( 2 \times 10^{15} \) cm\(^2\) s\(^{-1}\) in the coronal hole where \( N = 10^8 \text{ cm}^{-3} \) and \( T = 1.5 \times 10^6 \text{ K} \). This applies for the viscous transport of momentum parallel to the field, the transport across the field being smaller by a factor \( (\Omega t_c)^2 \), where \( \Omega = eB/Mc \) is the ion cyclotron frequency (10\(^9\) rad s\(^{-1}\) in 10 G) and \( t_c \) is the ion collision times, of about \( 10^7 \) s. Hence the effective viscosity is reduced to about \( 1 \text{ cm}^2 \text{ s}^{-1} \) for directions across the field, well below the resistive diffusion coefficient. Unfortunately the ion-ion collision time of about \( 10^8 \) s is comparable to the wave periods that might be effectively damped by viscosity along the field, so the classical viscosity provides only a rough estimate of the actual effect (see...

One wonders if there might be some mechanism that effectively converts fast-mode waves in the corona into slow-mode waves. As already noted, nonlinear effects (wave-wave interactions) are negligible because of the small amplitude of the Alfvén waves, while the interaction of Alfvén waves with the curvature of the field is weak because of the low $\beta$ of the coronal plasma. The other difficulty with conversion to slow-mode waves arises from the fact that the phase velocity $C$ of the Alfvén and fast-mode waves in the low-$\beta$ plasma is so much larger than the sound speed $c$. Oblique waves interacting with oblique boundaries fail to produce significant Fourier components with phase velocities as small as the speed of sound. So there is no evident way in which significant conversion to slow-mode waves can be effected within the corona.

There is an interesting coincidence, however, that should not go unnoticed, and that is the near equality of the sound speed $c = 2 \times 10^7$ cm s$^{-1}$ in the coronal hole and the Alfvén speed of about $1.4 \times 10^7$ cm s$^{-1}$ in the chromosphere below, where the number density is $2 \times 10^{10}$ cm$^{-3}$ (providing the same gas pressure at $7 \times 10^3$ K as $10^6$ cm$^{-3}$ provides at $1.5 \times 10^6$ K). If the transition region were made up of vertical fingers of chromospheric gas extending into the coronal gas for distances of the order of a wavelength or more ($\lambda = \alpha \approx 2 \times 10^8$ cm for $\tau = 10^2$ s), the Alfvén waves in the intruding chromospheric fingers would transfer a significant fraction of their energy to slow-mode waves in the corona, with relatively little reflection of the Alfvén waves. The slow-mode waves would be dissipated in a few wavelengths by Landau damping. On the other hand, observation appears to rule out tongues of chromospheric gas extending $10^7$ cm or more up into the corona, and the small scale height ($\approx 200$ km) of the cool dense chromospheric gas would appear to rule them out on theoretical grounds.

Hollweg (1984, 1986a; Sterling & Hollweg 1984) has pointed out that effective heating of the active X-ray corona by dissipation of Alfvén waves requires a damping effect comparable to the dissipation of individual eddies in a turbulent cascade at large Reynolds number. An eddy of scale $l$ and characteristic velocity $v(t)$ has an effective coherence time $\tau_l = l/v(t)$ during which it is dismembered and its energy transferred to smaller eddies with characteristic scale $l' = \alpha l$, where $\alpha$ is, by Kolmogorov's hypothesis, independent of $l$. Hollweg's "Kolmogorov" hypothesis is simply that a wave must somehow be made to disintegrate in a period of the order of $\tau_l$ (the wave period) if the wave is to contribute significantly to heating the active corona.

Conditions are not quite so demanding in the coronal hole. The dissipation must occur in a characteristic length of about one solar radius, instead of the dimensions ($10^9$--$10^{10}$ cm) of an active region. An Alfvén wave with a period of $10^5$ s propagates one solar radius in about 10 wave periods, rather than one wave period. But even so we are at a loss to disintegrate the wave so quickly.

In summary, then, there is no theoretical scheme of which we are aware by which an upward flux of Alfvén waves with characteristic periods of 30--300 s can be significantly dissipated within the first one or two solar radii. Indeed, as described in the previous section, it is not evident how an Alfvén wave flux as large as $5 \times 10^{16}$ ergs cm$^{-2}$ s$^{-1}$ can be introduced into the corona. As already noted, studies of line widths by Athay & White (1979a) and Leer et al. (1982) indicate an upper limit of $1 \times 10^{15}$ ergs cm$^{-2}$ s$^{-1}$ on Alfvén waves entering the corona, from observational results that are consistent with zero (see also Cheng, Doschek, & Feldman 1979).

We expect that the actual Alfvén wave flux is in fact close to the inferred upper limit because of the energy ($1 \times 10^{15}$ ergs cm$^{-2}$ s$^{-1}$) necessarily deposited in the far coronal hole to provide the observed high-speed solar wind. As already noted, Habbal & Leer (1982) have shown that such waves are dissipated at radial distances of 5--10 $R_\odot$.

To take another view of this situation, it follows that if Alfvén waves are not significantly dissipated in the first 1--2 $R_\odot$, then, on the basis of conservation of energy, the Alfvén wave flux into the base of the coronal hole cannot exceed the $1 \times 10^{15}$ ergs cm$^{-2}$ s$^{-1}$ necessary for the final acceleration of the solar wind to 400--600 km s$^{-1}$. For if $5 \times 10^{15}$ ergs cm$^{-2}$ s$^{-1}$ were present in the outer coronal hole, conservation of energy requires that it show up somehow in the solar wind, either as Alfvén waves, or as kinetic energy of the solar wind if the waves are dissipated. But this would require something of the order of 3 ergs cm$^{-2}$ s$^{-1}$ at the orbit of Earth. Wave amplitudes of $10^{-3}$ G in the ambient interplanetary field of $10^{-4}$ G would be required, or else average wind velocities of $10^5$ cm s$^{-1}$, neither of which are observed. The observed fluctuations of $10^{-4}$ G or less and wind speeds of only 400--600 km s$^{-1}$ put a theoretical upper limit of the order of $1 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ on the upward flux of Alfvén waves in the coronal hole.

4. ALTERNATIVES

If waves fail to provide the basic heat input to the lower end of the coronal hole, required by Withbroe's (1988) analysis, then we must cast about for alternatives.

Spicules obviously stir and heat the corona and have been investigated as a coronal heating agent by several authors. The conclusion (cf. Rabin & Moore 1980; Athay & Holzer 1982; Hollweg 1982b; Sterling & Hollweg 1984) is that the upward energy flux within the spicule may be as large as $10^6$ ergs cm$^{-2}$ s$^{-1}$, but the spicules occupy no more than about $10^{-2}$ of the total area, so that the mean is only of the order of $10^5$ ergs cm$^{-2}$ s$^{-1}$.

It is interesting to note that our present quest for alternatives to wave heating follows a well-trodden path. On quite general terms Tucker (1973) and Levine (1974) pointed out many years ago that the estimates of available wave amplitudes, and the dissipation of those waves, were too small to account for coronal heating. They were primarily concerned with the active X-ray corona, of course, but their remarks are not inappropriate in the present context. They suggested instead that the principal heat source comes from the dissipation of quasi-static fields, which can be accomplished only if those fields contain internal thin current sheets to provide enhanced dissipation. That is to say, there must be current sheets across which the change $AB$ in the field occurs over a transverse scale $\Delta l$ so small as to provide the necessary heat input (see eq. [2]). If one assumes the conventional resistivity $\eta \approx 3 \times 10^9$ cm$^2$ s$^{-1}$,
and a $\Delta B$ extending over a characteristic height $h = 10^{10}$ cm, the energy ($\Delta B^2/8\pi$) must be dissipated in a time $t_p$, where

$$h(t/h)\Delta B^2/8\pi \approx I \text{ ergs cm}^{-2} \text{ s}^{-1},$$

(6)

to produce an energy input $I$ ergs cm$^{-2}$ s$^{-1}$. If $\Delta B = 1$ G, and $I = 1 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$, it turns out that $t_p = 10^5$ s. If there is one such current sheet every l cm, the characteristic diffusion velocity $\eta/\Delta l$ must equal $I/t_p$, so that $\Delta l = \eta I t_p/l$. Putting $l = 10^7$ cm yields $I/t_p = 10^4$ cm$^{-2}$ s$^{-1}$ and $\Delta l = 0.3$ cm. As a matter of fact, the electron conduction velocity $u$ becomes equal to the ion thermal velocity of $2 \times 10^5$ cm s$^{-1}$ in the coronal hole when $\Delta l$ declines to $2 \times 10^5$ cm (for $\Delta B = 1$ G). Hence $\Delta l$ can hardly be less than $2 \times 10^5$ cm. For if $\Delta l$ were to be smaller, the increase in the electron conduction velocity would drive the ion-acoustic instability to produce much stronger plasma turbulence, resulting in a soaring anomalous resistivity that could drive the effective value of $\eta$ up to $10^{10}$ cm$^2$ s$^{-1}$ or more (cf. Friedman & Hamburger 1969; Smith & Priest 1972, 1973; Wentzel 1978). But the anomalous resistivity needs only to reach $2 \times 10^5$ cm$^{-2}$ s$^{-1}$ to provide $I/t_p = 10^4$ cm$^{-2}$ s$^{-1}$ with $\Delta l = 2 \times 10^5$ cm.

The theoretical development of the idea was carried out only in the abstract at that time (Wentzel 1978; Golub et al. 1980) because it was not understood how such current sheets are created (see also the recent work of Strauss 1988). Attention was directed principally to the magnetic loops of the active corona, for which it was clear that current sheets are essential for creating the X-ray corona. Subsequently it was pointed out (Glenross 1975, 1980; Parker 1981b, 1983b, 1990c) that the basic magnetostatic theorem, that almost all field topologies produce internal current sheets as an intrinsic part of static equilibrium (Parker 1972, 1979, 1981a, 1982, 1983a, c, d, 1986a, b, 1987a, 1989a, b, c, 1990a, b), automatically provides the necessary current sheets in the bipolar magnetic fields of active regions. On this basis, the active X-ray corona is to be regarded as a continuing flurry of nanoflares (Parker 1988; see also Sturrock & Uchida 1981).

It is clear, however, that the same explanation cannot be made to work for the coronal hole, for the simple reason, already mentioned, that the transverse components of the field are free to propagate away to infinity along the field. On the other hand, it must be appreciated that the energy input to the coronal hole is only a small fraction—about $1/20$—of the energy input to the active X-ray corona ($5 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ as compared to $10^7$ ergs cm$^{-2}$ s$^{-1}$). Hence, heating effects that are neglected as insignificant in the theory of the active corona must be reexamined in the context of the coronal hole. In particular, we must give serious attention to the small-scale magnetic fields that appear in the supergranule boundaries—network fields of some hundreds of gauss—and the flux bundles that pop up in the updrafts in the central regions of the supergranules—the intranetwork fields of some tens of gauss—and are swept into the boundaries where they interact with the network fields (Martin 1984, 1988; Porter et al. 1987; Wang et al. 1990) in association with vigorous microflaring. The rate of flux emergence has been studied by Liggett & Zirin (1985).

Martin (1984) was the first to argue that these small-scale fields are the major player in coronal heating. Porter & Moore (1988) have elaborated the idea further. The development in the preceding sections of the present paper indicates that there is no known alternative to the idea that the small scale fields—the network and intranetwork fields—are the principal source of heat for the coronal hole. It brings us again to the idea that the principal heat input is the result of dissipation of magnetic energy at thin current sheets. Porter and Moore suggest that the microflares, which account for a major part of the energy release, arise where small-scale ($2 \times 10^3$ km) magnetic bipole are jostled together by subphotospheric convection. Martin points out that most of the activity arises where newly emerged intranetwork fields are swept in among the established network fields. The general idea is simply that the small-scale fields, made up of unipolar and bipolar flux components, possess in miniature the major active features of the large "normal" active regions. The small bipole undoubtedly develop internal tangential discontinuities, just like their larger counterparts, where magnetic dissipation heats the gas trapped in the bipole. The small bipole reconnect where they are pushed against other bipole, or are pushed against unipolar fields. There may be no coronal mass ejections, no escaping fast particles, and no miniature sunspots. But at a low level, magnetic energy is continually increased by the emergence of new flux bundles and by the continual deformation of the bipolar fields by the subphotospheric convection. This magnetic energy of the small-scale fields is converted into heat at thin current sheets, in the manner with which we are familiar in the "normal" active regions (Parker 1981b, 1983b, 1987c, 1988, Porter & Moore (1988) estimate an average overall magnetic dissipation rate of about $5 \times 10^8$ ergs cm$^{-2}$ s$^{-1}$ in coronal holes as a consequence of the frequent microflares, of about $10^{26}$ ergs per flare. Most of this energy is injected into the near coronal hole, in the form of jets of gas, lashing flux bundles, and superheated plasma from a bipole interior suddenly freed by a reconnection into the ambient unipolar field. Some fraction of the energy release goes into Alfvén waves, because the reconnection releases field lines from one configuration to swing at the Alfvén speed into equilibrium in another configuration. Presumably it is the Alfvén waves generated in the corona by the activity of the small-scale fields that propagate outward to heat the distant coronal hole.

5. THE SMALL-SCALE MAGNETIC FIELDS

Consider some of the basic facts and figures concerning the small-scale fields. The network fields consist of flux bundles occasionally as large as $10^{25}$ maxwells, extending down to current limits of detection at about $10^{17}$ maxwells. The flux bundles congregate in the converging flows around the downdrafts at the network junctions where three supergranules meet (Zwaan 1981). The individual flux bundles are of mixed sign and are fairly closely packed in the vicinity of the network junctions. For the most part, then, the observed magnetic flux bundles represent the footprints of small bipole, with dimensions of $2 \times 10^5$ km or more. Smaller bipole are generally not resolved by existing magnetographs, so they do not appear as magnetic reversals, nor do they extend much above the chromosphere into the corona. The characteristic field strengths associated with the individual flux bundles in the network are of the order of $10^5$ G (Machado & Moore 1986; Martin 1984, 1988, 1990; Porter et al. 1987; Porter & Moore 1988; Wang et al. 1990), Fiedler & Cally (1990) and Simon et al. (1983), among others, have studied the vertical structure of individual unipolar flux bundles.

The intranetwork flux bundles are not so concentrated, with fields of the order of $10^{-3}$ G, as they first appear in the central updrafts of the supergranules (Martin 1988). They become increasingly concentrated as they are swept into the supergranule boundaries and move toward the downdrafts at the

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the junctions of boundaries. Much of the activity is caused by the intrusion of the intranetwork flux bundles into the established clusters of network fields. Wang et al. (1990) make the interesting point that the continual sweeping of the intranetwork flux bundles into the active network regions is the opposite of the popular idea about the spreading of field by turbulent diffusion of fields at the surface of the Sun. The point seems to be that insofar as the flux bundles diffuse across the surface of the Sun, it must be the supergranules, rather than the granules, that make the major contribution, with the local effect controlled to some degree by unseen convective motions well below the surface.

Porter & Moore (1988) note that the transient and spotty heat input of the microflares is consistent with the strong coronal density inhomogeneities inferred from the coronal emission spectrum (Orrall et al. 1990; Fiedler & Cally 1990) and is consistent with the small-scale velocities of 25 km s\(^{-1}\) inferred from line widths in the spectrum of the transition region.

Roumeliotis (1991) has constructed a simple theoretical model of the intense current sheets between colliding bipoles. He finds that the theoretical radiation spectrum has the same general form of differential emission measure as is observed. Fiedler & Cally (1990) obtain comparable results in a somewhat different model by introducing turbulent heat transport.

Now the large-scale magnetic field of the coronal hole is more or less unipolar with a strength of 5–10 G. The field converges into the network where it enters the Sun. The network region, occupying perhaps one-tenth of the surface of the Sun, is populated by both the magnetic flux bundles that make up the general magnetic field of the hole and by the flux bundles that form the many network bipoles. Hence all of the coronal hole communicates along the magnetic field into the active forest of bipoles. The suggestion is simply that the lower end of the coronal hole is heated principally by the magnetic activity of the network bipoles emanating upward along the mean field into the coronal hole.

The construction of a quantitative picture of the network fields, in terms of small bipoles and the field of the coronal hole (see sketch in Porter & Moore 1988) depends upon the relative proportions of flux bundles of different sign at the observed surface. The majority sign determines the sign of the unipolar field in the coronal hole above. The flux with the minority sign closes locally to form bipolar forms with flux bundles of opposite sign. One may imagine that the magnetic connections continually change as the crowded flux bundles are shoved around by the convection. Insofar as a flux bundle has a long-term identity at the photosphere, one may find it connected into a bipole partnership with another photospheric flux bundle at one time or another, while intermittently connected into the general field of the coronal hole between bipole partnerships. In this way the plasma in the coronal hole is heated by the magnetic dissipation of the continuing reconnection and by direct injection of heated plasma from within a bipole. There is not sufficient information to provide a quantitative assessment of the different contributions, but the following estimates spell out the requirements.

6. ENERGY CONSIDERATIONS

Denote by \(B_f\) the field intensity in the individual magnetic fibrils where they are observed, but generally unresolved, as they pass through some level in the solar atmosphere. Let \(\epsilon\) represent the fraction of the total area occupied by \(B_f\). If the network occupies one-twentieth of the solar surface, and the flux bundles occupy one-fifth of the area within the network, the result is \(\epsilon = 10^{-2}\). The individual flux bundles have one sign or the other, so denote by \(\alpha\) the fraction of the total flux that has the dominant sign. It follows that the fraction with the minority sign is \(1 - \alpha\), and the net flux, which has the dominant sign, is the fraction \(\alpha - (1 - \alpha) = 2\alpha - 1\) of the total.

If the mean unipolar field across the local portion of the coronal holes is \(B\), it follows that

\[
B = \epsilon(2\alpha - 1)B_f.
\]

As an example, then, suppose that \(B = 10 G\), with \(\epsilon = 10^{-2}\) appropriate for a coronal hole, and suppose that \(\alpha = \frac{1}{2}\). The result is \(B_f = 200B = 2 \times 10^5 G\), which is not out of line with other estimates.

The magnetic fluxes in the individual fibrils or bundles are detected down to a flux \(\Phi \approx 10^{17}\) maxwells (cf. Martin 1984, 1988, 1990; Wang et al. 1990). In general, a fibril radius \(R\) leads to the relation \(\Phi = \pi R^2 B_f\). If \(B_f\) is as strong as \(2 \times 10^5\) G, this corresponds to a fibril radius of 40 km for \(10^{17}\) maxwells. A flux of \(10^{18}\) maxwells corresponds to 120 km. Studies of the vertical structure of magnetic fibrils (Gabriel 1976; Withbroe & Noyes 1977; Simon et al. 1983; Deinzer et al. 1984a, b; Hasan 1985; Fiedler & Cally 1990) shows them to spread out rapidly with height above the photosphere, with a scale height \(h\) of the order of 800 km, as the ambient gas pressure declines, so that the fibril fields expand to fill the entire space within \(\sim 2 \times 10^3\) km above the surface. Very roughly, then, the energy of a magnetic fibril is of the order of \(\pi R^2 h (B_f^2/8\pi)\), with an average energy \(\delta = eh(B_f^2/8\pi)\) per unit surface area. With \(h = 800\) km, \(B_f = 2 \times 10^5\) G, and \(\epsilon = 10^{-2}\), this leads to \(\delta \approx 10^{11}\) ergs cm\(^{-2}\) s\(^{-1}\). If this energy is dissipated at an average rate \(I = 5 \times 10^5\) ergs cm\(^{-2}\) s\(^{-1}\), the characteristic dissipation time is \(\delta/I = 2 \times 10^5\) s. That is to say, if the forest of bipoles is reconnected once a day, there is plenty of energy released to heat the coronal hole. Some of the energy may remain trapped within the bipoles and be radiated away from the chromosphere and transition region. However, the active reconnection of the fields, producing the microflares described by Porter & Moore (1988) and unresolved nanoflares, introduces a major fraction of the dissipated magnetic energy into the surrounding corona.

As already noted, some fraction of the magnetic energy release goes into Alfven waves. It was pointed out in § 2 that the wave flux of \(1 \times 10^5\) ergs cm\(^{-2}\) s\(^{-1}\) to heat the distant coronal hole involves an rms wave velocity in any one transverse direction of about 12 km s\(^{-1}\). Note that the characteristic wave periods produced by the network microflares are not limited to the photospheric periods of 50–300 s. Rapid reconnection between small intense bipoles may produce waves with periods as short as 1 s or less. It is not possible at present to estimate what fraction of the energy of the network microflares is subject to immediate loss by radiation and what fraction is transmitted outward into the coronal hole by thermal conduction, high-speed jets, slow- and fast-mode magnetohydrodynamic waves, and Alfven waves.

Now if the field is dissipated by random reconnections in a period of \(10^5\) s, and the characteristic fibril radius is, say, 100 km, the mean cutting rate is \(10^5\) cm s\(^{-1}\). This is a modest rate indeed, when we note the observational examples (Glackin 1975; Sheeley et al. 1975a, b; Howard & Svestka 1977; Rust & Webb 1977; Wolfson et al. 1977; Nolette et al. 1977; Heyvaerts, Priest, & Rust 1977) or reconnection of entire fields over dis-
tances of $10^6$ cm or more in 10–20 hr, implying characteristic cutting speeds in excess of $2 \times 10^4$ cm s$^{-1}$. In this way the newly emerging complex fields of active regions relax toward potential forms in a matter of hours.

7. CONCLUSION

An examination of existing theory of the generation of waves in the subphotospheric convection, and of their transmission up a coronal hole, indicates that it is only Alfven waves that can make the trip, and that they are not generated sufficiently strongly as to supply more than a fraction of the necessary $5 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ in the coronal hole. An upper limit of about $1 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ has been inferred from observations. What is more, there is no known theoretical mechanism to explain the necessary dissipation of such waves in the first $1$–$2 R_\odot$ to account for the principal heat input to the coronal hole. Alfven waves of $10^7$ s period may dissipate effectively over distances of $5$–$20 R_\odot$, but not in $1$–$2 R_\odot$.

The only available heat source of sufficient strength is the activity of the network magnetic fields, which are continually replenished and activated by the arriving intranetwork fields. The network microflares alone, observed in association with the small-scale fields, indicate that magnetic energy is released at the necessary $5 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$. The agitation caused by the microflares, and perhaps less conspicuous reconnection events, may be presumed to excite slow, fast, and Alfven wave modes in the corona, as well as high-speed jets and small blast waves. One presumes that $1 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ or more goes into Alfven waves or fast-mode waves with periods of the order of $10^7$ s to heat the outer coronal hole and provide the high-speed wind streams. The rest of the energy is evidently supplied to the inner coronal hole by the modes that dissipate more quickly (i.e., in $10^5$ s).

It should be noted that the active network, that is, the network in active regions, probably supplies more heat to the corona than in coronal holes. But the heat budget of the active corona, at $10^7$ ergs cm$^{-2}$ s$^{-1}$, is 20 times larger than in the coronal hole, so the effect of the network activity is presumably minor in an active region. By the same token, the network activity may be the major supplier to the quiet corona, whose energy requirement is estimated to be only $5 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ by Withbroe & Noyes (1977). So little is known of conditions in the quiet corona that no more can be said.

Now the small-scale activity of the network fields presumably involves a variety of forms of reconnection. Bipoles reconnect with other bipoles where the photospheric convection happens to ram them together. Bipoles reconnect with unipolar flux bundles where they are pressed together. The general swirling and twisting of the bipoles produce internal tangential discontinuities, whose dissipation may greatly heat the plasma trapped with each bipole, and such superheated plasma is free to escape into the coronal hole when the bipole reconnects with a unipolar field. The reconstructions, which cannot be observed effectively, permit some peculiar walking around of the footpoints of the mean unipolar field of the coronal hole, as well as some “spontaneous” disappearance of flux when the individual flux bundles are not resolved in the magnetograph. Some of these possibilities are outlined in another paper (Parker 1991).

The essential point is that future progress in understanding the details of the small-scale magnetic activity depends in large part on observation, which currently is working on the problem, but it is hampered by the limited spatial and temporal resolution. Adaptive optics on the ground and a large space telescope dedicated to observing the Sun are essential to an eventual full understanding of the small-scale activity. The activity is of fundamental importance in its own right and as the primary architect of the heliosphere.

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