CHROMOSPHERIC AND CORONAL HEATING MECHANISMS

U. NARAIN
Institut für Theoretische Astrophysik, Universität Heidelberg,
Im Neuenheimer Feld 561 D–6900 Heidelberg, Federal Republic of Germany

and

P. ULMSCHNEIDER
Astrophysics Research Group, Physics Department,
Meerut College Meerut - 250001, India

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Abstract. We review the mechanisms which have been proposed for the heating of stellar chromospheres and coronae. These consist of heating by acoustic waves, by slow and fast mhd waves, by body and surface Alfvén waves, by current or magnetic field dissipation, by microflare heating and by heating due to bulk flows and magnetic flux emergence. Some relevant observational evidence has also been discussed.

1. Introduction

Ground based and satellite observations have shown that probably all stars with the possible exception of A-stars have shells or regions in their outer atmosphere in which the temperature is much higher than the photospheric value. These chromospheric and coronal layers are characterized by large energy loss: chromospheres lose energy predominantly by radiation and coronae by conduction, radiation and stellar wind flows. To prevent these hot layers from rapidly cooling down to the photospheric boundary temperature a considerable amount of mechanical heating must constantly be applied. To affect this necessary heating many different mechanisms have been proposed. Reviews of these chromospheric or coronal heating processes were given among others by Van de Hulst (1953), Billings (1966), Osterbrock (1961), Schatzman and Souffrin (1967), Kuperus (1969), Piddington (1973), Bray and Loughhead (1974), Stein and Leibacher (1974, 1980), Athay (1976), Withbroe and Noyes (1977), Wentzel (1978, 1981), Mewe (1979), Ulmschneider (1979, 1981), Leibacher and Stein (1980, 1982), Chiuderi (1981, 1983), Hollweg (1981a, 1983, 1985b), Kuperus et al. (1981), Priest (1982, 1983), Golub (1983), Kuperus (1983), Heyvaerts (1984, 1985), Cassinelli and MacGregor (1986), Parker (1983c, 1986), Rosner et al. (1986), Ulmschneider (1986), Kumar (1987). For reviews on flare-related heating which we do not discuss in this work see Priest (1981, 1982) and Spicer (1982).

In recent years it became clear, however, that very likely the heating phenomena in chromospheres and coronae can not be explained by a single process but are rather due to the action of a multitude of mechanisms. Some of these may operate globally, others only in particular physical situations. At the present time it is not possible to decide which mechanisms are the important ones, as in spite of a considerable effort the observational basis for each proposed process is rather thin and the poor state of theoretical development of many mechanisms could only lead

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to an unfair verdict. In this work we thus attempt to give a comprehensive review of the chromospheric and coronal heating mechanisms which have been proposed to date and try to be as unbiased as possible. As we are interested in physically defined mechanisms we do not discuss the literature on assumed heating mechanisms where the damping length is treated as a free parameter.

For the purpose of presenting the abundant and sometimes disparate literature it was necessary to adopt a classification scheme of the various mechanisms for which we took the type of mechanical energy input into the heated layers as a guide. This guide could, however, not always be followed strictly. Sections 2 to 5 (acoustic wave heating, heating by fast and slow mhd waves, Alfvén body waves, Alfvén surface waves) deal with waves which in late-type stars are thought to be excited by rapid turbulent and convective motions at the base of the photosphere. In other cases it is possible that waves may be generated by small-scale instabilities, or by mode-coupling from other waves. Section 6 discusses current (magnetic field) dissipation. Here the energy is introduced into the coronal magnetic loops by slow convective motions and by the differential rotation of the photospheric foot points or by buoyancy. As it gives a unifying picture, the resonant LCR circuit approach is treated in a separate Section 7. Section 8 describes the heating by microflares. The heating by bulk flows or by magnetic flux emergence is discussed in Section 9. In Section 10 some relevant observational evidence is presented. Section 11 contains our conclusions.

2. Acoustic wave heating

Only a few years after Edlén's (1941) discovery that the solar corona is a gas layer with a million degree temperature, a theoretical explanation for this hot layer, which attributes the heating to the dissipation of acoustic shock waves was offered by Biermann (1946, 1948) and by Schwarzschild (1948). Subsequently this idea was developed further by Schatzman (1949), De Jager (1959), De Jager and Kuperus (1961), Osterbrock (1961), Bird (1964a,b, 1965), Kuperus (1965, 1969, 1972), Kopp (1968), Stefanik (1969), Stein (1969), Stein and Schwartz (1972), Gonczi et al. (1977), Ulmschneider (1967, 1970, 1971a, 1974, 1976, 1989).

The basic idea is that stellar surface convection zones, as is the case for all turbulent flow fields, invariably generate a spectrum of acoustic waves. The mechanical energy flux F_M carried by the upwardly propagating acoustic waves for sufficiently high frequency is given by

$$F_M = \rho \, \overline{v^2} \, c_S \tag{2.1}$$

where ρ is the density, v the velocity and c_S the sound speed. As the temperature in the photospheric layers does not vary much and if radiation damping is small the conservation of wave energy flux requires that

$$v \sim \rho^{-1/2}. \tag{2.2}$$

The density decrease of the outer atmosphere thus results in a rapid growth of the wave amplitude and due to the nonlinear terms in the hydrodynamic equations leads to shock formation, shock dissipation and heating of the outer stellar layers.

In the following subsections we discuss first the shock wave heating theory. Other acoustic wave heating processes which do not involve shocks are the heating by radiation damping and by ionization pumping. The latter two processes are discussed in Sections 2.6 and 2.7.

2.1. Weak Shock Theory

For linear sawtooth waves with pressure and velocity variations $p = p_0 + p_m - 2p_m t/P$, $v = v_m - 2v_m t/P$ where P is the wave period, t the time and subscript m indicates maximum amplitudes the wave energy flux (*erg cm*$^{-2}$ *s*$^{-1}$) is given by:

$$F_M = \frac{1}{P}\int_0^P (p - p_0)\; v\; dt = \frac{1}{3}\; p_m v_m \approx \frac{1}{12}\;\gamma\; p_0\; c_S\; \eta^2 \quad , \tag{2.3}$$

where p_0 is the unperturbed pressure, γ is the ratio of specific heats and where for weak shocks one has for the total jumps $2p_m \approx \gamma p_0 \eta$, $2v_m \approx c_S \eta$. Here the shock strength $\eta = (\rho_2 - \rho_1)/\rho_1$, where ρ_1, ρ_2 are the densities in front and behind the shock (Ulmschneider 1970). The shock dissipation rate (*erg cm*$^{-3}$ *s*$^{-1}$) of the wave can be written

$$\epsilon_M = \frac{\rho T \Delta S}{P} = \frac{\rho_0\; c_S^2}{\gamma(\gamma-1)P}\ln\left(\frac{p_2}{p_1}\left(\frac{\rho_2}{\rho_1}\right)^{-\gamma}\right) \approx \frac{1}{12}\;\frac{\gamma(\gamma+1)}{P}\; p_0\; \eta^3 \quad , \tag{2.4}$$

where ΔS is the entropy jump per unit mass at the shock front. Assuming in analogy to ray optics that the quantity $F_M c_S^2$ is conserved, differentiating with respect to height z and identifying the result with Eq. (2.4) gives

$$\frac{d\eta}{dz} = \frac{\eta}{2}\left(\frac{\gamma g}{c_S^2} - \frac{3}{2c_S^2}\frac{dc_S^2}{dz} - \frac{(\gamma+1)\;\eta}{c_S P}\right) \quad , \tag{2.5}$$

where g is the gravitational acceleration and γ is assumed constant. The refractive term involving dc_S^2/dz is small in the chromosphere but may become large in the transition layer. Eq. (2.5) shows that irrespective of the initial shock strength the shocks eventually reach a *limiting shock strength*

$$\eta^{lim} = \frac{\gamma\; g\; P}{(\gamma+1)\; c_S} \tag{2.6}$$

and a limiting wave flux

$$F_M^{lim} = \frac{1}{12}\;\frac{\gamma^3\; g^2\; P^2}{(\gamma+1)^2\; c_S}\; p_0 \quad . \tag{2.7}$$

Eq. (2.5) can also be written

$$\frac{dF_M}{dz} = -\frac{1}{c_S^2}\frac{dc_S^2}{dz}\; F_M - \frac{F_M}{L_A} \quad , \tag{2.8}$$

where L_A is the acoustic damping length given by

$$L_A = \left(\frac{\gamma\, p_0\, c_S^3}{12\,(\gamma+1)^2}\right)^{1/2} F_M^{-1/2} P = \frac{c_S\, P}{(\gamma+1)\,\eta} \quad . \tag{2.9}$$

From Eq. (2.6) it is seen that upon reaching the limiting shock strength the damping length becomes equal to the scale height $L_A^{lim} = c_S^2/\gamma g$. Heating laws of the type (2.5) or (2.8) assume that the shocks have small amplitude such that the weak shock relations are satisfied and that radiation damping from the waves is small which close to the stellar surface can be a bad approximation (Ulmschneider 1988).

The validity of Eq. (2.6) for the prediction of the limiting shock strength has recently been investigated by Cuntz and Ulmschneider (1988) using nonlinear time-dependent wave calculations. They find that for short period waves in non- ionizing, isothermal atmospheres with constant gravity the value predicted by Eq. (2.6) is closely reached. The basic property of acoustic shock waves to reach a limiting strength is also maintained in non-isothermal, ionizing atmospheres with height dependent gravity. Eq. (2.6) predicts the limiting strength in these more realistic atmospheres only if the actual values of the sound speed and γ are used and if short period waves ($P < P_A/5$, c.f. Eq. 2.10) are considered. For longer period waves Eq. (2.6) increasingly underestimates the limiting strength.

The weak acoustic shock theory was used to compute coronal models of late-type stars (Couturier et al. 1980, Hammer 1982) and of early-type stars (Hearn and Vardavas 1981). The propagation of weak acoustic shock waves guided by diverging magnetic fields through a static model of the solar chromospheric network and transition layer was studied by Foukal and Smart (1981).

2.2. Strong shock and time-dependent treatments

The restriction to weak shocks can be overcome by adopting the principle of shape similarity invariance found experimentally. In this approach based on the work of Brinkley and Kirkwood (1947) and Bird (1964a,b), the amount of shock heating depends on the way by which the post shock state relaxes (e.g. Ulmschneider 1967, Flower 1977, Gonczi et al. 1977, Souffrin 1982, see also Bray and Loughhead 1974). In the Weyman cycle rapid radiative relaxation is followed by adiabatic expansion, in the Schatzman cycle adiabatic expansion is followed by slow radiative relaxation. For the solar chromosphere the Weyman cycle appears valid (Ulmschneider et al. 1978). Using the strong shock approach of Bird (1964b) coronal hole models were computed by Pineau des Forêts (1979) and by Flower and Pineau des Forêts (1983).

As there is a sensitive mutual interaction between the radiative losses and the shape of the wave, more detailed treatments describe the development of acoustic waves with time-dependent radiation hydrodynamic methods. In these treatments the shocks may have arbitrary strength. For a review of time-dependent calculations of radiative (magneto-) hydrodynamic wave propagation in stellar atmospheres see Ulmschneider and Muchmore (1986). Klein et al. (1976) investigated Lyman continuum emitting transients in A-star atmospheres. Hydrodynamic wave propagation along coronal loops was discussed e.g. by Mariska and Boris (1983) or McClymont and Canfield (1983). For treatments of shock propagation with radiation pressure for early-type stars see Section 2.8.

Calculations of the short period (compared to P_A, c.f. Eq. 2.10) acoustic wave propagation in the solar chromosphere (Ulmschneider et al. 1978, Schmitz et al. 1985, Ulmschneider et al. 1987) using various methods of radiation treatments show that acoustic heating can in principle explain the temperature minimum (as region of shock formation) and the magnitude of the chromospheric temperature rise. Moreover shock dissipation is well capable to balance the empirically determined chromospheric radiation losses. Direct viscous or thermal conductive heating is many orders of magnitude too small.

The long period acoustic waves, the 5 min- and 3 min- oscillations (see Sec. 10), which are an outstanding observational effect on the sun are not candidates for solar chromospheric heating because they are standing waves and the observed phase shift of 90^o between velocity and temperature fluctuations (Deubner 1974, Fleck and Deubner 1989) precludes that these waves form shocks in the chromosphere. Long period acoustic waves other than the short period waves are associated with nonradial vibrations of the star. In late- type giant stars this acoustic wave type eventually leads to large scale non-radial oscillations and pulsations, which in Mira stars are observed to produce extensive shock heating in the outer layers (Hinkle et al. 1982). Long period acoustic waves are also thought to be instrumental for generating mass loss in late-type giant stars (Cuntz 1989).

Using short period acoustic waves, chromospheric models for late-type stars other than the sun were constructed by Ulmschneider et al. (1979), Schmitz and Ulmschneider (1980a, 1980b, 1981) as well as Ulmschneider (1989). As in the solar case it was found that in principle the acoustic heating appears to be a viable mechanism. Temperature minima are explained as heights where shock formation occurs or as upper boundaries of radiation damping zones (Schmitz and Ulmschneider 1981). These time-dependent calculations show that the acoustic shock waves always approach a limiting shock strength roughly predicted by Eq. (2.6). The calculations suggest moreover that short-period waves are dissipated low in the atmosphere while the dissipation of longer-period waves occurs at greater heights (e.g. Stein and Leibacher 1974, 1980, Ulmschneider 1979).

Sterling and Hollweg (1988) extend and develop the rebound shock model of Hollweg (1982b) for spicules. The authors study adiabatic acoustic wave propagation and shock formation in a specified rigid flux-tube geometry. They find that the model is capable of generating structures with properties consistent with observations (except for the temperature and the life-time).

2.3. Acoustic energy generation in late-type stars

Renzini et al. (1977), Fontaine et al. (1981) and others used the Lighthill (1952, 1954) formula for quadrupole sound generation to estimate the acoustic energy production on the basis of theoretical convection zone models for late-type stars. These formulae were developed for isotropic turbulence in non- gravitational atmospheres. Stein (1967, 1968) showed that in addition to quadrupole contributions there are contributions of monopole and dipole sound generation. He found for the sun that under reasonable assumptions about the turbulent energy spectra of the convection the generated acoustic spectrum peaks at a period of around $P = 30s$ which is roughly one tenth of the acoustic cut-off period, P_A, at the top of the convection zone,

$$P_A = \frac{4\pi c_S}{\gamma g} \quad . \tag{2.10}$$

Dimensional arguments (Stein 1981, Ulmschneider and Stein 1982) show that the acoustic flux generated in the convection zones of late type stars with effective temperature T_{eff} by the monopole, dipole and quadrupole source terms is given by

$$F_M^{mono}/\sigma T_{eff}^4 \sim g^{-0.19}\, T_{eff}^{2.1} \;, \tag{2.11}$$

$$F_M^{dip}/\sigma T_{eff}^4 \sim g^{-0.58}\, T_{eff}^{6.4} \;, \tag{2.12}$$

$$F_M^{quad}/\sigma T_{eff}^4 \sim g^{-0.96}\, T_{eff}^{10.6} \;, \tag{2.13}$$

where σT_{eff}^4 is the radiative flux of the star. Bohn (1981, 1984) employed Stein's (1967) method using improved opacities and a better treatment of molecules in his convection zone models to calculate acoustic energy fluxes and acoustic spectra for a large sample of late-type stars. For stars where the convection zone is not too thin, his results can be approximated by

$$F_M = 1.4 \cdot 10^{26}\, T_{eff}^{9.75} g^{-0.5} \alpha^{2.8} \;, \tag{2.14}$$

where α is the ratio of mixing-length to pressure scale height. Bohn moreover found that the acoustic spectra rise from the cut-off period P_A to a peak near $P_A/10$ with a subsequent rapid decrease towards shorter periods. For late-type main sequence stars the peak of the acoustic spectra occurs at longer wave periods. The above computations of the acoustic wave generation have recently been criticized (see Ulmschneider 1989), but we feel that the acoustic spectrum generated by the turbulent surface convection zones is probably not greatly changed, and lies roughly within a factor of ten of the acoustic cut-off frequency of the top of the convection zone.

A basic property of the acoustic energy generation in the convection zones of late-type stars is that the acoustic flux depends only on the parameters which specify such zones, (T_{eff}, g, α), and thus is *independent of rotation.* In addition it must be stressed that the general tendency of acoustic waves to develop into limiting strength shock waves (outside of radiation damping zones) results in minimum acoustic fluxes given by Eq. (2.7). If the acoustic spectrum is not changed much by more realistic computations and similar as in Bohn's calculations shows a peak near $P_A/5$ or $P_A/10$, then Eqs. (2.7), (2.10) allow a limiting acoustic wave flux estimate. Ulmschneider (1989) finds

$$F_M \approx 2.3 \cdot 10^4\, p \approx 2.3 \cdot 10^4\, g\, m \quad , \tag{2.15}$$

if $P = P_A/10$ or a four times higher flux if $P = P_A/5$. Here p is the gas pressure and m is the mass column density. Note that this wave flux results from the steepening property of the atmosphere and is *independent of the generation process by the convection zone.* Such a limiting acoustic shock wave behaviour also applies for longitudinal magnetohydrodynamic tube waves in situations where the magnetic tube is fully spread and thus has constant cross section with height.

2.4. Acoustic energy generation in late-type stars

The discovery that stars with the same T_{eff} and gravity show greatly different UV and X-ray emission (Basri and Linsky 1979, Vaiana et al. 1981) and that this activity is related to rotation (see e.g. Rutten 1987) demonstrated that chromospheric emission in typical stars is not due to acoustic heating but is rather connected with a magnetic heating mechanism. Recently Oranje and Zwaan (1985) as well as Schrijver (1987a, 1987c) argued that there are *two basic contributions* to the chromospheric heating of late-type stars: the acoustic heating which is independent of rotation and constitutes a minimum contribution for a given T_{eff} and the rotation related magnetic heating which is usually a much larger contribution. This would explain why the most slowly rotating stars for a given T_{eff}, the so called basal flux stars, show an intrinsic lower limit of chromospheric emission.

In addition, the large acoustic contribution due to the steep T_{eff}- dependence of Eq. (2.14) could explain the low range of activity observed for early F-stars (Walter and Schrijver 1987). Judge (1989) finds that the chromospheric emission from non-Mira M and C giants depends only on T_{eff} and attributes this to acoustic heating, see also Querci and Querci (1985). The fact that late-type giant stars for large B-V do not show chromospheric emission variability but have only a T_{eff}- dependent Ca II emission as found by Middelkoop (1982, Fig. 4c) also points to pure acoustic heating. That the large gravity dependence of the acoustic energy generation (c.f. Eq. 2.14) does not imply a large gravity dependence of the chromospheric emission has been shown by Ulmschneider (1988, 1989). This could explain the observed fact that the basal chromospheric emission flux does not depend much on gravity (Rutten 1987).

2.5. Radiative heating

So far in Sections 2.1 to 2.5 we have discussed the acoustic shock wave heating theories for late-type stars. The so called heating by radiation damping is another process by which a large amount of acoustic wave energy can be dissipated. Ulmschneider (1971b), Ulmschneider et al. (1978), Schmitz and Ulmschneider (1981), Schmitz et al. (1985) and Ulmschneider (1988) showed that radiation damping of acoustic waves due to the H^- emission produces extensive radiation damping zones in the deeper photospheric layers of stars. Giovanelli (1979) and Ulmschneider et al. (1987) showed that radiation damping by lines is also an important contributor to the chromospheric wave dissipation. Radiation damping, however, does not lead to enhanced mean temperatures. On the contrary, large amplitude waves due to the nonlinearity of the Planck function produce a depression of the average temperature (Ulmschneider et al. 1978). Kalkofen et al. (1984) showed that in spite of this depression the chromospheric line emission is enhanced, which by an outside observer may be interpreted as an apparent temperature enhancement.

2.6. Heating by ionization pumping

Lindsey (1981) notes that slow adiabatic compression of a partially ionized gas wherein ionization equilibrium is maintained will show a response which approaches that of an un- ionized gas under isothermal compression

$$\frac{\delta p}{p} \approx -\frac{\delta V}{V} \quad , \tag{2.16}$$

where p is the gas pressure and V the volume. After Lindsey the same change of state is accomplished by a rapid adiabatic compression followed by an isobaric irreversible ionization relaxation. The adiabatic compression part of the hysteresis loop is described by

$$\frac{\delta p_1}{p} = -\gamma \frac{\delta V_1}{V} \quad , \tag{2.17}$$

where γ is the ratio of specific heats. With $\delta V \approx \delta V_1 + \delta V_2$ and $\delta p \approx \delta p_1$ for the total volume and pressure changes in the hysteresis loop the volume change of the irreversible part, δV_2 is given by

$$\delta V_2 \approx -\Big(\frac{\gamma - 1}{\gamma}\Big) \frac{V}{p}\, \delta p \quad . \tag{2.18}$$

The heat ΔQ (erg) accumulated per cycle is

$$\Delta Q = -\int p\, dV = -\frac{1}{2}\, \delta p\, \delta V_2 = \frac{1}{5}\frac{V}{p}\, (\delta p)^2 \quad , \tag{2.19}$$

where the integral runs along the closed hysteresis loop and where $\gamma = 5/3$ has been used. As the relaxation process occurs with a relaxation time t_R, heating occurs only if $P < t_R$ where P is the wave period. For high frequency waves satisfying this condition a heating rate ($erg\ cm^{-3}\ s^{-1}$)

$$\epsilon_{IP} = \frac{\Delta Q}{V\, t_R} = \frac{(\delta p)^2}{5\, p}\, \frac{1}{t_R} \quad , \tag{2.20}$$

is found while low frequency waves do not show heating by this process.

According to Lindsey, who calls this heating process ionization pumping, the relaxation times t_R range from $\sim 20\ s$ at the base of the hydrogen plateau ($h \sim 1000\ km$) to $\sim 90\ s$ in the upper regions ($1600\ km \leq h \leq 2200\ km$) for high frequency oscillations. For lower frequencies, he believes the dissipation to be less efficient. He concludes that the irreversible relaxation to the ionization equilibrium during and following compression of the chromospheric medium can serve as an efficient mechanism for dissipating compressional energy of hydrodynamic waves into heat. Unlike shock dissipation, ionization remains

3. Heating by fast and slow magnetoacoustic body waves

3.1. The wave modes

In the absence of gravity and for a perfectly conducting, compressible, uniform gas in a *homogeneous* magnetic field, there exist three types of body waves, the *Alfvén-mode, fast-mode* and *slow-mode waves.* The latter two wave types are also called magnetoacoustic waves as these waves show gaspressure fluctuations. Alfvén mode waves which will be discussed in Sec. 4, to first order do not show gaspressure or density fluctuations. For the derivation of these waves from the linearized set of magnetohydrodynamic equations see e.g. Jeffrey (1966), Bray and Loughhead (1974), Priest (1982). The phase speed of Alfvén waves is

$$v_{ph,A} = v_A \ \cos\theta \quad , \tag{3.1}$$

where θ is the angle between the wave propagation vector $\mathbf{k}$ and the magnetic field $\mathbf{B}_0$ and v_A is the Alfvén speed given by

$$v_A = \frac{B_0}{(4\pi\rho_0)^{1/2}} \quad . \tag{3.2}$$

The phase speeds v_F, v_S of the fast-mode and slow-mode waves are given by

$$v_F^2 = \frac{1}{2}\Big((c_S^2 + v_A^2) + \Big[(c_S^2 + v_A^2)^2 - c_S^2\ v_A^2\ \cos^2\theta\Big]^{1/2}\Big) \quad . \tag{3.3}$$

$$v_S^2 = \frac{1}{2}\Big((c_S^2 + v_A^2) - \Big[(c_S^2 + v_A^2)^2 - c_S^2\ v_A^2\ \cos^2\theta\Big]^{1/2}\Big) \quad . \tag{3.4}$$

About ten years ago it was found that the magnetic field at the solar surface is not homogeneous but appears rather in the form of highly concentrated flux-tubes which are located at the boundaries of the granulation and supergranulation cells (Stenflo 1978, Zwaan 1978). A powerful approach to treat these flux tubes is to assume that the physical quantities do not change much over the tube cross-section and that they are given essentially by their values at the tube axis. For these so called thin flux-tubes there exist three wave modes (Spruit 1982): the *longitudinal or sausage mode*, the *transverse or kink mode* and the *torsional mode.* The latter two wave types to first order do not show fluctuations of the tube cross-section and thus do not have gaspressure variations. These waves thus correspond to the Alfvén waves in homogeneous fields and will be discussed in Sec. 4.

Longitudinal waves have axisymmetric variations of the tube cross-section and show gaspressure fluctuations. Actually, the main restoring force of longitudinal tube waves is the gas pressure. Herbold et al. (1985) found that longitudinal tube waves are essentially acoustic waves which propagate along the magnetic flux-tube. Defouw (1976), Roberts and Webb (1978) and others showed that these waves propagate with the tube speed c_T:

$$c_T = \left(\frac{c_S^2\ v_A^2}{c_S^2 + v_A^2}\right)^{1/2} \quad . \tag{3.5}$$

For cases where $c_S << v_A$ (e.g. in the high chromosphere) the longitudinal tube waves thus correspond to the slow-mode waves, while there is no analog of the fast-mode waves in the thin flux-tube situation.

3.2. Characteristics of the Wave Modes

The material motion in longitudinal tube waves is along the magnetic field, while for transverse and torsional waves it is perpendicular to the field. Due to horizontal pressure balance the magnetic flux-tubes spread rapidly with height and at the so called canopy height (somewhere in the middle chromosphere) fill out the entire available space (Jones and Giovanelli 1982, Anzer and Galloway 1983). It appears that the longitudinal, transverse and torsional waves which propagate along the magnetic flux-tubes at photospheric and low chromospheric heights go over into the slow-, fast- and Alfvén- wave modes in the locally homogeneous fields above the middle chromosphere. Considerable nonlinear coupling between the wave modes is expected to occur at these canopy heights.

The fast-mode waves can propagate in any direction. For weak magnetic fields where the sound speed $c_S >> v_A$ one has $v_F \approx c_S$, while, for strong fields, $c_S << v_A$ one finds $v_F \approx v_A$. In both cases the propagation velocity is independent of the direction, while at intermediate fields the velocity depends weakly on the direction. In weak fields the direction of material motion is longitudinal and the waves are essentially acoustic waves, while in strong fields the motion is perpendicular to the direction of the field, and the waves are essentially magnetohydrodynamic in character.

Since fast-mode waves can carry energy across magnetic fields the energy requirements of distinct coronal features can be met by waves of modest energy flux density entering the corona over a large surface area. During propagation they are refracted into regions of low Alfvén velocity (Habbal et al. 1979) hence they can deposit their energy at certain preferred sites even though the energy might have entered the atmosphere uniformly from below. For a given observed rms velocity fluctuation the energy flux density that can be carried by the fast-mode waves is roughly v_A/c_S times the corresponding quantity for the sound waves.

For the heating of the corona one most disadvantageous characteristic of fast-mode waves is the possibility that these waves approaching the transition region from below are heavily refracted and can be totally internally reflected and thus may never reach the corona. The Brewster angle is only 6^o if the ratio of Alfvén speeds above and below the transition region is 10 (Hollweg 1981a). On the other hand fast- and slow- mode waves in the corona may be generated by mode-coupling from other wave types, see Sec. 4.

The slow-mode waves can propagate only in directions close to the direction of the magnetic field, and both at weak and strong fields, the only allowed direction of propagation is exactly along the field, while at an intermediate field ($c_S = v_A$), the allowed directions lie within a cone with half angle 27^o. The propagation velocity of slow-mode disturbances, along the direction of the field, is approximately v_A for $c_S >> v_A$ and c_S for $c_S << v_A$. The material motion is perpendicular to the direction of propagation for weak fields and is in the direction of the field for strong fields.

3.3. Weak shocks

Consider linear sawtooth waves similar to those in Sec. 2.1 which propagate in the x-direction. The undisturbed magnetic field B_0 is assumed to be inclined by the angle θ with respect to the x- direction such that it has the components B_{0x} and B_{0y}. The wave energy flux can then be written (Osterbrock 1961):

$$F_M = \frac{1}{P} \int_0^P \left[\left((p - p_0) + \frac{1}{8\pi}(B^2 - B_0^2) \right) v_x - \frac{1}{4\pi} v_y B_{0x} \left(B_y - B_{0y} \right) \right] dt \quad , \tag{3.6}$$

where for fast-mode waves and in the weak shock approximation one has

$$F_M^{fast} = \frac{1}{12} \rho_0 \, v_F \left(c_S^2 + v_A^2 \sin^2\theta / (1 - v_A^2 \cos^2\theta / v_F^2)^2 \right) \eta^2 \quad . \tag{3.7}$$

where the shock strength $\eta = (\rho_2 - \rho_1)/\rho_1$ is defined as in Sec. 2.1. In a similar way to Sec. 2.1 the shock dissipation rate ($erg\ cm^{-3}\ s^{-1}$) for fast-mode waves (Osterbrock 1961) is given by

$$\epsilon_F \approx \frac{\rho_0}{4\ P} \left(\frac{1}{3}(\gamma + 1)\ c_S^2 + v_A^2 \sin^2\theta / (1 - v_A^2 \cos^2\theta / v_F^2)^2 \right) \eta^3 \quad . \tag{3.8}$$

Eqs. (3.7) and (3.8) allow us to derive the fast-mode damping length L_F

$$L_F = \frac{F_M^{fast}}{\epsilon_F} = \frac{v_F\ P}{3\ \eta} \left(\frac{c_S^2 + v_A^2 \sin^2\theta / (1 - v_A^2 \cos^2\theta / v_F^2)^2}{c_S^2\ (\gamma + 1)/3 + v_A^2 \sin^2\theta / (1 - v_A^2 \cos^2\theta / v_F^2)^2} \right) \approx \frac{v_F\ P}{3\ \eta} \quad . \tag{3.9}$$

For slow-mode waves similar equations apply. Osterbrock (1961) states that slow-mode disturbances can be discussed as pure gasdynamic or acoustic shocks, propagating only along the direction of the magnetic field, and they can be analyzed by omitting the magnetic field and refraction effects. The fluxes, dissipation rates and damping lengths for slow-mode waves are given by

$$F_M^{slow} \approx \frac{1}{12}\ \rho_0\ c_S^3\ \eta^2 \quad , \tag{3.10}$$

which is identical to Eq. (2.3) and

$$L_S = \frac{F_M^{slow}}{\epsilon_S} \approx \frac{c_S\ P}{3\ \eta} \quad . \tag{3.11}$$

which is roughly equal to Eq. (2.9).

If shock dissipation by magnetoacoustic waves is to be an important heating mechanism for chromospheres or coronae it is essential that the conditions of shock formation are favourable. In the nonlinear process of shock formation the high velocity parts of the wave profile catch up with the low velocity parts. With a velocity amplitude v, shocks form after a time $t_{sh} = \lambda / 4v$, where λ is the wavelength. For fast-mode waves this occurs at a distance $d_{sh}^{fast} = t_{sh}\ v_F = v_F^2 P/4v$. This shows that for a given wave energy flux short period waves form shocks much more readily. Because $d_{sh}^{fast}/d_{sh}^{slow} = (v_F/c_S)^{5/2} \approx 316$ for $v_A = 10\ c_S$ it is seen that it is much more difficult for fast-mode waves to form shocks than for slow-mode or acoustic waves. Shock dissipation thus is a likely process if one has a short wave period (compared to P_A, c.f. Eq. 2.10) and a low propagation speed (c_S, c_T). For the sun this amounts to $P < 60\ s$ and propagation speeds $< 10\ km/s$.

3.4. Heating in a plane-parallel medium

Osterbrock (1961) provided the first systematic study of the generation, propagation and dissipation of magnetoacoustic (and Alfvén) waves in a plane-parallel solar atmosphere. His study is limited to weak shock waves as discussed above, with a wavelength much smaller than the scale height. Based on the now discarded picture of a homogeneous magnetic field distribution he proposes that the shock dissipation of fast -mode waves is the dominant heating mechanism in chromospheres. Because of the outward density decrease, the Alfvén speed increases much faster than the sound speed. This leads to much more pronounced refraction of fast-mode waves than sound waves, so less fast- mode energy is likely to reach coronal layers. The weak shock theory of magnetoacoustic waves was used by Uchida (1963) to construct solar corona models.

For application to the solar wind Barnes (1966, 1968) develops a general theory of collisionless plasma heating by mhd waves. He finds that the damping of fast-mode waves can heat both protons and electrons, enhancing the component of the kinetic temperature parallel to the mean magnetic field. D'Angelo (1968, 1969) suggested that Landau damping of ion acoustic waves results in heating of the corona. This has been studied in greater detail by Revathy (1977).

Mäckle (1969) studies the propagation of fast-mode shock waves in an atmosphere permeated by a nearly horizontal, uniform magnetic field. The shock structure has been computed in detail. He finds that the waves may reach the upper chromosphere and may heat it.

Musielak (1982, 1987) computes theoretical models of homogeneous chromospheres for main sequence stars assuming heating by acoustic, magneto-acoustic and Alfvén waves. In his later work he finds that energy carried by fast-mode (or acoustic) waves is significantly dissipated in deep photospheric layers, mainly because of radiative damping. He points out that magnetically heated coronae might be formed around DA and DB stars provided a mechanism for energy dissipation by transverse modes exists.

Schwartz and Bel (1984) investigate the propagation of magnetoacoustic gravity modes in the solar atmosphere with a magnetic field inclined at an arbitrary angle to the direction of gravity. In magnetoacoustic gravity waves simultaneously three restoring forces, magnetic tension, gas pressure and buoyancy are considered. According to them strong magnetic fields yield essentially acoustic modes propagating along field lines whereas weak magnetic fields produce a complicated competition amongst the various forces and are inefficient in transporting energy to coronal heights. As the other modes are presumed to be dissipated low in the chromosphere the authors conclude that only Alfvén-mode waves seem to remain as a viable option for coronal heating.

Hollweg (1985b) estimated the heating by fast-mode waves in active region loops and coronal holes. From Eq. (3.8) with $v_F \approx v_A \gg c_S$ and $\eta \approx v'/v_A$ (Osterbrock 1961) one finds a fast-mode shock heating rate ($erg\ cm^{-3}\ s^{-1}$)

$$\epsilon_F \approx \frac{B_0^2}{16\ \pi\ P} \left(\frac{v'}{v_A}\right)^3 \quad , \tag{3.12}$$

where v' is the velocity amplitude. For active region loops with $B_0 = 100\ G$ and $v_A = 3000\ km/s$ Hollweg finds $\epsilon_F = 6 \cdot 10^{-4}\ erg\ cm^{-3}\ s^{-1}$ for waves with $P = 100\ s$ and $v' = 200\ km/s$. For a coronal hole with $B_0 = 5\ G$ and $v_A = 800\ km/s$ he obtains for the same wave $\epsilon_F = 7.8 \cdot 10^{-5}\ erg\ cm^{-3}\ s^{-1}$. These heating rates appear adequate. However, taking for the assumed v' observed velocity amplitudes in the inner corona (Withbroe 1988), which are by a factor of ten smaller, then these heating rates are reduced by a factor of 1000 and no longer adequate, except if one would suppose a wave period of $P \approx 1\ s$.

3.5. Heating in a structured medium

Using geometrical optics techniques, Milkey (1970) shows that fast-mode waves generated in a non-homogeneous, anisotropic medium can heat the regions of strong magnetic field in the chromosphere above the boundaries of supergranular cells via shock dissipation. He offers his result as an explanation for the origin of the Ca^+ emission network. Durrant and Michalitsanos (1971) do not agree with this contention because it does not explain the enhancement of the photospheric emission network.

Habbal et al. (1979) propose that coronal loops may be heated via collisionless damping of fast-mode waves. They represent loop-like field structures by a dipole field with the center of the dipole at a given distance below the solar surface and determine the density of the medium by hydrostatic equilibrium along the field lines under isothermal conditions. They consider waves with a period of 2 s. This choice has been made to keep the wavelengths small enough so that ray-tracing techniques may be used. Also it allows reasonable damping lengths in the corona. It should be noted however that waves with a period of 2 s would be damped by Joule-viscous frictional damping in the chromosphere and would be unable to reach the corona from the convection zone (Hollweg 1981a, Kuperus et al. 1981, Heyvaerts 1985). In order to avoid this difficulty they postulate the existence of fast-modes of such short periods at the base of the corona.

Following Habbal et al. the group velocity of fast-mode waves under coronal conditions is given by

$$\mathbf{v}_g = \frac{dbfx}{dt} = \frac{\partial \omega}{\partial \mathbf{k}} = v_A \hat{k} \quad , \tag{3.13}$$

where $\hat{k}$ is the unit vector in propagation direction. This shows that the waves can transport energy across field lines. Further they have

$$\frac{d\mathbf{k}}{dt} = -\frac{\partial \omega}{\partial bfx} = -k \, \frac{\partial v_A}{\partial bfx} \quad , \tag{3.14}$$

which shows that waves propagating outward from the coronal base are refracted into regions of low Alfvén speed. The waves are assumed to damp by Landau/transit time damping (Barnes 1966). This form of damping is collisionless and assumes $\omega t_{pp} \geq 1$ where t_{pp} is the proton- proton collision time. With $t_{pp} = 17\, T^{3/2}/(\ln \Lambda \, n_p)$ given by Braginskii (1965), where n_p is the proton density and $\ln \Lambda \approx 21$ the Coulomb logarithm and with values $T = 1.8 \cdot 10^6 \; K$, $n_p = 1 \cdot 10^8 - 5 \cdot 10^8 \; cm^{-3}$ for quiet region loops (see Table I) we find $\omega t_{pp} = 6 - 30$. This shows that the collisionless condition is satisfied for the $P = 2\pi/\omega = 2 \; s$ waves which Habbal et al. consider. When $\omega t_{pp} \leq 1$, the fast-modes still damp but by collisional processes such as Joule, viscous and heat conduction damping (Hollweg 1981a) or by shock dissipation.

The interesting feature of Landau damping is that it increases with the plasma β ($= 8\pi p/B^2$) if β is small. This occurs because the number of resonant particles increases with β. In the isothermal model of Habbal et al., β is largest where v_A is smallest. However, the fast-mode waves refract toward regions of lower v_A and thus to areas of larger damping. This means that these waves can lead to anisotropic heating in the corona. The effect is enhanced by a " positive feedback" process in which local heating increases β which in turn increases the heating rate as the number of resonant particle increases in the heated regions. This may lead to a situation in which once fast-mode waves begin heating the plasma, they dump all their energy in a small volume of space. Habbal et al. suggest that such processes can account for the highly structured nature of the active region corona in a natural way. They even suggest that cool cores are a natural consequence of coronal heating by fast-mode waves in accordance with the following picture: Fast-mode waves impinge on a cylindrical structure from the outside, dump all their energy on the outside of the cylinder (being unable to penetrate to the interior) and thus resulting in a cool core. Habbal et al. conclude that fast-mode waves can heat loops provided the wave energy flux density at the coronal base is of the order $10^5 erg \; cm^{-2} \; s^{-1}$.

Hollweg (1981a) points out that the general level of the observed coronal and transition region non-thermal line widths is consistent with a flux of fast-mode waves that is large enough to provide the coronal energy requirements. But it needs to be demonstrated that the non-thermal velocities correspond to fast-mode waves. He feels that in order to make refinements to theoretical calculations of Habbal et al. more detailed observational knowledge of how $\mathbf{B}_0, \beta$ and v_A vary (both inside and outside loops) in the corona is required. Hollweg argues that because fast-mode waves refract into regions of low Alfvén velocity these waves can not be important for the heating of active regions which are clearly more energetic than other regions. This argument, however, is not safe as it does not consider the possible generation of fast-mode waves by mode-coupling.

According to Hollweg (1981a) and Heyvaerts (1985) in a low β -plasma where fast-mode waves obey the dispersion relation $\omega^2 = k^2 v_A^2$, the propagation vector $\mathbf{k}$ can be split into a horizontal (conserved) component k_H and a vertical component k_V. One has

$$k_V^2 = \frac{\omega^2}{v_A^2} - k_H^2 \quad . \tag{3.15}$$

If $k_H > \omega/v_A$, the waves are evanescent and carry no energy upwards. Based on the observed $\omega - k_H$ diagram corresponding to photospheric and chromospheric motions (Hollweg 1978b) it is conjectured that the fast-mode waves in the corona should be evanescent. This conjecture may be false if the corona is structured on scales comparable to or smaller than the fast-mode wavelength. It should be noted that active regions are characterized by larger values of v_A and presumably larger values of k_H than the quiet corona and coronal holes. These trends work in the direction of enhancing the evanescent region of the $\omega - k_H$ diagram.

Zweibel (1980) studies the thermal stability of coronal fine structure, loops and streamers heated by fast-mode waves. According to Zweibel, when the wave period P is short compared to the electron and ion collision times, t_e and t_i, the waves are damped by collisionless wave-particle interaction (Landau damping) where wave energy is exchanged with particles moving at the resonant speed along the magnetic field. When $t_e < P < t_i$, damping by electron thermal conduction and collisionless ion damping takes place, while in the case $t_i < P$ ion-viscous damping occurs.

Habbal and Leer (1982) study the dissipation of fast-mode mhd waves in the solar wind. They find that these waves, propagating outwards from the sun in coronal hole regions, dissipate primarily through collisionless interaction with electrons rather than protons. They find that in the wind accelerating region this dissipation can lead to higher electron than proton temperatures.

Rae and Roberts (1982) study the propagation of mhd waves in inhomogeneous plasmas. They point out that the cusp resonance excited by propagating slow-mode waves may be important for heating the solar corona.

Leroy and Schwartz (1982) investigate magnetoacoustic gravity modes in an isothermal exponential atmosphere permeated by a uniform vertical magnetic field. They find the inhomogeneous analogues of the fast- and slow-mode waves in a homogeneous medium. The computations of Schwartz and Leroy (1982) show that magnetoacoustic gravity waves are not a major factor in coronal energetics because coronal fast-mode waves are evanescent (see Hollweg 1981a) and slow-mode waves are mainly acoustic in nature.

Cargill and Priest (1983) study heating of postflare loops by slow mhd shocks. They conclude that slow shock heating is a viable method for heating post-flare loops.

Flå et al. (1984) examine the possibility that fast-mode waves transport energy from magnetically closed regions into coronal holes and deposit most of it in the region of supersonic flow of high-speed solar wind streams. Their study indicates a broad range of coronal hole parameters for which fast-mode waves can play such a role. They conclude that improved knowledge of the large scale coronal magnetic structure is required to confirm the results of their study.

Herbold et al. (1985) study the propagation of non-linear radiatively damped longitudinal tube waves along magnetic flux- tubes in the solar photosphere and chromosphere. They find that these waves readily transport mechanical energy to chromospheric heights and dissipate it via shock formation very similar to acoustic waves. They conclude that longitudinal tube waves are good candidates for chromospheric heating.

Merzljakov and Ruderman (1986) study non-linear slow body wave propagation in a magnetic slab embedded in a magnetic environment. They conclude that the five-minute oscillations can generate body waves that may heat the upper corona.

Ulmschneider (1986) shows that owing to acoustic shock formation longitudinal tube waves seem to satisfy best the observational constraints which a chromospheric heating mechanism must satisfy: the rapid onset of chromospheric emission above the solar temperature minimum with a high net radiative cooling rate reaching values of 0.3 $erg\ cm^{-3}s^{-1}$, the magnetic-field related spatially inhomogeneous solar activity, the rotation related stellar chromospheric activity with its low gravity dependence and the emission gap near A-stars.

3.6. Wave energy generation

Stein (1981) studies the dependence of the energy flux of mhd waves, generated by turbulent motions in stellar convection zones, on the stars effective temperature, surface-gravity and magnetic field strength. He concludes that Alfvén and acoustic slow-mode waves generated in a background magnetic field sufficiently strong to dominate the gas motions could be prime candidates for heating stellar chromospheres and coronae.

Taking equipartition values for the magnetic field $B^2 = 8\pi p$ (Saar 1987) and estimating the gas pressure p at the top of the convection zone of late-type stars Ulmschneider and Stein (1982) as well as Ulmschneider (1986) estimated the mhd wave fluxes generated in the convection zone. They find

$$F_M^{fast}/\sigma T_{eff}^4 \simeq 1.2 \cdot 10^{-38}\ g^{-0.959}\ T_{eff}^{10.6}\ , \tag{3.16}$$

$$F_M^{slow}/\sigma T_{eff}^4 \simeq 9.4 \cdot 10^{-9}\ g^{-0.192}\ T_{eff}^{2.13}\ , \tag{3.17}$$

$$F_M^{Alfven}/\sigma T_{eff}^4 \simeq 8.6 \cdot 10^{-9}\ g^{-0.192}\ T_{eff}^{2.13}\ . \tag{3.18}$$

Musielak and Rosner (1987, 1988) study the generation of mhd waves by turbulent motions in a gravitationally stratified and magnetized fluid. They find monopole, dipole and quadrupole emissions to be responsible for the generation of the slow- and fast-mode waves. The Alfvén mode can be generated by dipole emission only. In their later paper they find that a uniform magnetic field does not lead to any significant increase in mhd wave generation. They conclude that this type of mhd wave generation cannot explain the observed stellar coronal emission.

Musielak et al. (1987, 1989) calculate the generation of longitudinal waves in magnetic flux tubes embedded in an otherwise magnetic field-free turbulent and compressible fluid. The waves are assumed to be excited by external turbulence. They are generated by dipole emission. The wave flux produced by this method, however, is found to be two orders of magnitude too small to play a significant role for the chromospheric heating.

In summarizing Section 3 we feel that the acoustic-like slow- mode waves or longitudinal tube waves with their easy shock formation must play an important role in the chromospheric and probably also in coronal heating. The low generation rate of these waves found by Musielak et al. (1987, 1989) may be an artifact of their excitation by small amplitude turbulent shaking. If the picture by Parker (1981c, Fig. 2) of large horizontal displacements and compressions of the magnetic flux tubes produced by the motion of convective bubbles applies, then much larger nonlinear effects play a role and much larger longitudinal and transverse wave fluxes are expected to be generated. In addition, nonlinear mode-coupling always seems to lead to longitudinal or slow-mode waves. Due to the ready shock dissipation these waves very likely are considerably damped in the low and middle chromosphere.

For fast-mode waves the situation is less clear. The strong rise of the Alfvén speed in the high chromosphere leads to strong refraction and to difficulties in forming shocks. However, even if the transmission of fast-mode waves from deeper layers into the high chromosphere and corona is inefficient, these waves will not be absent at higher layers, due to the production by mode- coupling from transverse and torsional Alfvén waves. This occurs e.g. in regions where the field lines are bent and where fields in which the latter waves propagate meet adjacent fields.

4. Heating by Alfvén body waves

4.1. Wave types, direct dissipation

The principal restoring force of Alfvén waves is magnetic tension. Alfvén waves propagate along the magnetic field with a phase speed given by Eq. (3.1). The group speed of these waves is equal to the Alfvén speed v_A, their energy always propagates along the direction of the magnetic field in qualitative agreement with observations indicating that the heating of the solar atmosphere is associated with the magnetic field. Alfvén waves are pure shear waves and to first order are not affected by the compressibility of the medium or by buoyancy forces. The waves never show evanescence or total internal reflection. In this section we consider *Alfvén body waves* in homogeneous magnetic fields as well as *transverse-* and *torsional body Alfvén waves* in magnetic flux tubes.

Alfvén (1947) was first to propose that the high temperature of the corona may be due to the dissipation of shear Alfvén waves, generated as a result of motions of granules in the photosphere. Since such waves are always associated with electric currents there will be Joule heating as the medium has a finite conductivity. He used the transverse electrical conductivity in his estimation and found the period of the waves responsible for heating the corona to be of order of a few minutes. Subsequent investigations by Cowling (1953), Piddington (1956) and others showed that Joule and viscous damping are inadequate to dissipate these waves.

The damping length of Alfvén waves due to *Joule and viscous heating* after Osterbrock (1961) is given by

$$L_{JV} = \frac{P^2\ v_A^3}{4\pi^2(c^2/4\pi\overline{\lambda} + \mu/\rho)} \quad , \tag{4.1}$$

where P is the wave period, $\overline{\lambda}$ is the electrical conductivity (s^{-1}), μ is the viscosity ($g\ cm^{-1}\ s^{-1}$) and c is the velocity of light. The damping length due to the *ion-neutral frictional heating* is given by

$$L_{IN} = \frac{P^2\ v_A}{4\pi^2}\ \frac{1+\chi}{\chi\ t_{in}} \quad , \tag{4.2}$$

where χ is the ratio of the mass density of neutral atoms to the mass density of ions and t_{in} is the mean collision time of an ion to lose its momentum in collisions with neutral atoms (Osterbrock 1961). The combined effect is given by the sum of the reciprocals, so the effective damping length is $L_{eff} = (1/L_{JV} + 1/L_{IN})^{-1}$.

For comparison with chromospheric radiative cooling rates spatially averaged dissipation rates ($erg\ cm^{-3}\ s^{-1}$)

$$\epsilon_{JV} = \frac{F_M^{Alfven} f}{L_{JV}} \quad , \qquad \epsilon_{IN} = \frac{F_M^{Alfven} f}{L_{IN}} \quad , \tag{4.3}$$

have to be considered, where F_M^{Alfven} is the Alfvén wave flux and f is the filling factor, i.e. the ratio of the area permeated by the magnetic field to the total area. With an estimate of the Alfvén flux from Eq. (3.18) which gives 13 percent of the total solar flux σT_{eff}^4 and with a filling factor of 0.01, one finds (Ulmschneider 1986) $F_M^{Alfven} f = 8 \cdot 10^7\ erg\ cm^{-2}\ s^{-1}$ as well as $\epsilon_{JV} = 5 \cdot 10^{-5}$, $\epsilon_{IN} = 1 \cdot 10^{-3}\ erg\ cm^{-3}\ s^{-1}$. In this estimate $\overline{\lambda} = 2 \cdot 10^{12} s^{-1}$, $\mu = 5 \cdot 10^{-4}\ g\ cm^{-1} s^{-1}$, $\rho = 2.3 \cdot 10^{-9} g/cm^3$, $v_A = 11\ km/s$, $\chi >> 1$, $t_{in} = 7.5 \cdot 10^{-4}\ s$ and a wave period $P = 40\ s$ were used. Note that here we use a corrected value of $t_{in} = (n_p Q_{pH} v_{th})^{-1} \approx 6.3 \cdot 10^9\ T^{-1/2} n_p^{-1}$ where $n_p = 1.1 \cdot 10^{11} cm^{-3}$ is the proton density, $T = 5700\ K$, and v_{th} the mean thermal velocity of the proton. For the ion-neutral cross-section Q_{pH} see Nowak and Ulmschneider (1977, Table II). In the high chromosphere using values from model F of Vernazza et al. (1981) at 2000 km height and $B = 20\ G$ we obtain for the same wave period and wave flux, $\epsilon_{JV} = 3 \cdot 10^{-6}$, $\epsilon_{IN} = 6 \cdot 10^{-5}\ erg\ cm^{-3}\ s^{-1}$ with $v_A = 94\ km/s$ and $t_{IN} = 9 \cdot 10^{-4}\ s$.

The dissipation rates $\epsilon_{JV}, \epsilon_{IN}$ are much too low to balance the empirical chromospheric radiative cooling rates of Vernazza et al. (1981) of about 0.3 $erg\ cm^{-3}\ s^{-1}$ in the low and $10^{-2}\ erg\ cm^{-3}\ s^{-1}$ in the high chromosphere, even if the unknown wave period would be considerably decreased. As these dissipation rates vary only slowly with height they do not reproduce the observed rapid onset of the chromospheric emission behind the temperature minimum found by Vernazza et al. (1981). Direct dissipation of Alfvén waves via the above processes thus appears unimportant for the solar chromosphere.

Osterbrock's (1961) result that Alfvén waves are strongly damped by Joule, viscous and ion-neutral frictional dissipation in the photosphere is due to the fact that he assumed a homogeneous magnetic field of 2 G in his investigation. Because of the now observed extremely inhomogeneous field distribution with field strengths of $\sim$ 1500 G as discussed in Sec. 3.1 this result is no longer true. In fact, the damping of Alfvén waves for fields larger than $10 - 20\ G$ is so small that it is a problem to explain how they give up their energy before propagating away, either into the solar wind along open field lines or back down to the photosphere along closed fields (Priest 1982). More promising Alfvén wave dissipation mechanisms by resonant heating, turbulent heating, nonlinear interaction, Landau damping and phase-mixing are discussed below in Sections 4.2 to 4.5.

4.2. Resonant and turbulent heating

Hollweg (1978a, 1981b) considers the propagation of torsional Alfvén waves in open vertically directed magnetic flux tubes. For these tubes realistic photospheric, chromospheric and coronal temperature and density distributions are assumed. No coupling to gravity occurs. Hollweg shows that in the absence of dissipation the twist velocities v obey

$$\left(v_A^2 \frac{\partial^2}{\partial l^2} + \omega^2\right) B_0^{1/2} v = 0 \quad , \tag{4.4}$$

$$i\omega B' = B_0^{1/2} \frac{\partial}{\partial l} \left(B_0^{1/2} v\right) \quad , \tag{4.5}$$

where l is the distance along some field line near the axis of the flux tube, B' the field perturbation and ω is the angular frequency. For constant B_0 these equations go over to those for Alfvén waves in homogeneous fields. Representing the tube by a sequence of piecewise exponential atmospheric layers the equations are solved assuming a monochromatic torsional excitation $\sim \exp(i\omega t)$ with velocity amplitude $v = 1\ km/s$ at the bottom of the tube.

The most interesting result of this analysis is the appearance of distinct resonances. Hollweg attributes these to the strong reflection of the upward propagating wave on the steep Alfvén speed gradient in the upper chromosphere and transition layer which almost results in a standing wave pattern. Only at certain resonance frequencies, where the propagating and reflected waves interfere constructively, does significant energy transport to infinity occur.

Schwartz et al. (1984) in a similar more detailed study using the same boundary conditions as Hollweg, likewise find distinct resonances in open flux tubes.

Zugzda and Locans (1982) studying the transmission coefficient of Alfvén waves propagating through a series of slabs dispute the reality of these resonances for open tubes and point out the similarity with the case of two adjacent isothermal atmospheres with Alfvén speeds v_{A1}, v_{A2}. In this case the transmission coefficient D varies smoothly from a small value at low wave frequencies to a maximum value of $D = 4 v_{A1} v_{A2} (v_{A1} + v_{A2})^{-2}$ at high frequencies. For tubes at the boundary of supergranules and above sunspots they find smoothly varying transmission coefficients.

Spruit (1981) derived general equations of motion for longitudinal and transverse Alfvénic tube waves in thin flux-tubes embedded in a nonmagnetic compressible fluid. He discusses the propagation of the transverse waves and finds that these waves could transport significant amount of energy into the chromosphere.

Alfvén wave propagation in closed coronal loops leads to *loop resonances* as found by Hollweg (1981b, 1984a,b) and by Zugzda and Locans (1982). These resonances arise from reflections at the steep Alfvén speed gradients at both ends of the loop. The resonant wave periods are roughly given by

TABLE I

Total length $l_{||}$, temperature T and number density n in coronal loops after Priest (1982). Radiative loss rates after Eq. (6.13) are added except at low temperature, where we used the original thin plasma value.

	Interconnecting	Quiet region	Active region
$l_{\|\|}$ (km)	$2 \cdot 10^4 - 7 \cdot 10^5$	$2 \cdot 10^4 - 7 \cdot 10^5$	$10^4 - 10^5$
T(K)	$2 - 3 \cdot 10^6$	$1.8 \cdot 10^6$	$10^4 - 2.5 \cdot 10^6$
n (cm^{-3})	$7 \cdot 10^8$	$0.2 - 1.0 \cdot 10^9$	$0.5 - 5.0 \cdot 10^9$
ϵ_R (erg cm^{-3} s^{-1})	$1.4 \cdot 10^{-5}$	$2 \cdot 10^{-6} - 4 \cdot 10^{-5}$	$7 \cdot 10^{-4} - 3 \cdot 10^{-3}$

$$P_{res} = \frac{2 l_{||}}{m\ v_A} \quad , \tag{4.6}$$

where $m = 1, 2, 3,..$ is an integer and $l_{||}$ is the length of the coronal part of the flux tube. A short loop might have $l_{||} = 2 \cdot 10^4\ km, B = 100\ G$ and a proton number density $n_p = 10^{10} cm^{-3}$ which together with Eq. (4.6) gives $P_{res} = 20\ s,\ 10\ s,\ 7\ s$ etc. Longer loops with smaller Alfvén speeds would have longer resonant periods. The authors find that these loop resonances act like "windows" whereby a large energy flux can pass almost unimpeded from the convection zone into the corona. Hollweg estimates that an energy flux of $1.5 \cdot 10^7 erg\ cm^{-2} s^{-1}$ into a short loop may be available for heating which is quite adequate. Outside the loop resonances short period ($P \leq 10\ min$) undamped waves are not able to carry significant energy flux, because of the destructive interference.

Priest (1982) summarized the properties of coronal loops. To estimate the importance of the loop heating mechanisms we we reproduce in the following Table I his Tab. 6.2 leaving out the post-flare and simple-flare loops.

Ionson (1982, 1984) shows that the process of *resonant heating* by Alfvén waves in coronal loops can be well understood by considering the analogy to an LCR resonant electric circuit. As the LCR circuit approach may be generalized to include non-wave heating processes we discuss Ionson's work in greater detail in Section 7.

Hollweg (1984a,b) investigates the propagation of torsional Alfvén waves in a medium consisting of three layers, each having a different Alfvén speed. He finds two types of resonances. When the central layer (i.e. the cavity) has the largest (or smallest) of the three Alfvén speeds as in coronal loops the ordinary loop resonance occurs. A second type of resonance happens when the Alfvén speed of the central layer is intermediate between the other two speeds which may be important for solar spicules. Further if the energy lost by heating greatly exceeds the energy lost by leakage out of the central layer then heating can be independent of the specific damping process, in agreement with Ionson (1982) c.f. Sec. 7. For resonant heating Hollweg (1984a,b, 1985b) gives a heating rate ($erg\ cm^{-3}\ s^{-1}$)

$$\epsilon_H \approx \frac{4\ \pi^3\ h_A\ v_{Ac}\ \Sigma\ m\ P_\omega}{l_{||}^3}\ F_M^{Alfven} \quad , \tag{4.7}$$

where h_A is the scale height of the Alfvén speed in the chromosphere-corona transition layer. P_ω is a spectrum weighting factor related to the input power spectrum and is normalized such that $\int P_\omega d\omega = 1$. The summation is carried over the different resonances m (c.f. Eq. 4.6). F_M^{Alfven} is the incident Alfvén wave flux given by

$$F_M^{Alfven} = 2\ \rho_p\ \delta v_p^2\ v_{Ap} \frac{B_c}{B_p} = \pi^{-1/2} \rho_p^{1/2}\ \delta v_p^2\ B_c \quad . \tag{4.8}$$

Here the indices c and p indicate coronal and photospheric values, respectively, B_c/B_p is an area factor due to the spreading of the loop and the factor 2 arises from two possible polarizations. If $h_A = 200\ km, v_{Ac} = 2000\ km/s, l_\| = 10^5\ km, F_M^{Alfven} = 1.6 \cdot 10^8\ erg/cm^2 s, m = 1$ and $P_\omega = 60/2\pi = 9.5\ s$, where the power spectrum was assumed to be flat up to $P = 60\ s$, Hollweg (1985b) finds $\epsilon_H \approx 7.5 \cdot 10^{-4}\ erg/cm^3 s$, which is sufficient (see Table I as well as Hollweg and Sterling 1984) to power active region loops.

Hollweg and Sterling (1984) carry out model fitting based on the resonant heating theory of Hollweg (1984b) and the coronal loop data of Golub et al. (1980). Using Eqs. (4.7) and (4.8) they find that for reasonable forms of the input power spectrum and reasonable values for the wave velocity amplitudes in the photosphere the resonant heating theory is fully consistent with the data. For the resonant heating theory see also Sec. 7.

Hollweg (1984b, 1985a,b, see also Heyvaerts and Priest 1983) suggested that torsional Alfvén waves may be dissipated by Helmholtz-Kelvin instabilities with a subsequent generation of a Kolmogoroff-type turbulent cascade, transverse to the tube direction, where after reaching small enough wave numbers the energy is dissipated by viscosity or electrical resistivity. Hollweg finds a *turbulent heating* rate

$$\epsilon_T = \frac{\rho\ < v^2 >^{3/2}}{l_{corr}} \quad , \tag{4.9}$$

where l_{corr} is a correlation length in the direction across the flux-tube. Assuming $l_{corr} \approx r$ where r is the tube radius and identifying v with the total observed horizontal nonthermal velocities, Hollweg (1985a) found that ϵ_T agrees well with the empirical chromospheric radiative cooling rates of Avrett (1981) between the heights 800 and 2000 km. However, as pointed out by Ulmschneider (1986) this heating rate, due to the small photospheric tube radius r does not agree with the empirical cooling rates at heights below 800 km. In addition, the turbulent heating rate ϵ_T does not explain the empirically determined rapid onset of chromospheric emission above the temperature minimum. For additional discussions of turbulent heating in see Section 6.4.

Schüssler (1984) suggested that torsional Alfvén waves in magnetic flux tubes might be excited by the intergranular cyclonic downdraft of the cool gas outside the tubes.

Mariska and Hollweg (1985) study some nonlinear aspects of Alfvénic pulses of twist propagating in coronal loops and the underlying chromosphere. They find that the pulses can result in significant motion of the transition region and underlying chromosphere. These motions do not resemble spicules.

4.3. Dissipation by nonlinear interaction (mode - coupling)

A very efficient heating process by Alfvén waves is that through mode - coupling they transfer energy to other modes which dissipate more readily. Alfvén waves may be dissipated by nonlinear interaction with either a nonuniform field or with other Alfvén waves (Kaburaki and Uchida 1971, Chin and Wentzel 1972, Uchida and Kaburaki 1974, Wentzel 1974, 1976, 1977, Priest 1982 p.221). It is found that when the magnetic field is weak ($v_A < c_S$) two Alfvén waves traveling in opposite directions along a magnetic field line can couple nonlinearly to give an acoustic wave which dissipates via shock formation rather easily. The dispersion relations for these waves may be written as (Sagdeev and Galeev 1969)

$$\omega_0 = v_A \, k_0, \quad \omega_1 = v_A \, k_1, \quad \omega_2 = c_S \, k_2, \quad , \tag{4.10}$$

where $\omega_0, \omega_1, \omega_2$ and k_0, k_1, k_2 are the frequencies and wave numbers of the two Alfvén waves and the generated acoustic wave, respectively. As a result of coupling of the two incident Alfvén waves energy conservation requires that the frequency of the acoustic wave is given by

$$\omega_2 = \omega_0 + \omega_1 \quad . \tag{4.11}$$

The wave numbers are related by

$$k_2 = k_0 - k_1 \quad . \tag{4.12}$$

The minus sign results from the conservation of momentum as the two Alfvén waves are propagating in opposite directions. From Eqs. (4.10) to (4.12) one obtains

$$\frac{\omega_1}{\omega_0} = \frac{c_S - v_A}{c_S + v_A} \tag{4.13}$$

and

$$\omega_2 = \frac{2\omega_0 \, c_S}{c_S + v_A} \quad . \tag{4.14}$$

In regions of strong magnetic field ($v_A > c_S$), one Alfvén wave (ω_0, k_0) may decay into another Alfvén wave (ω_1, k_1) travelling in the opposite direction together with a sound wave (ω_2, k_2) travelling in the same direction. The interaction takes place provided the conditions of conservation of energy and momentum are satisfied. The resulting frequencies of Alfvén and sound waves are given by Eq. (4.13) with the RHS multiplied by -1 and (4.14). Obviously the resulting Alfvén wave has a smaller frequency than the original one. This process of decay of one Alfvén wave into another Alfvén wave of lower frequency plus an acoustic wave may continue till all the Alfvén energy is converted to acoustic waves and to thermal energy via shock dissipation of acoustic waves.

Wentzel (1974, 1976, 1977) finds a significant production of acoustic waves from Alfvén wave interaction in closed coronal loops if v_A/c_S lies between 1/30 and 30. The dissipating acoustic waves preferentially heat the top of the loop. He concludes that in this manner short period ($P \leq 100\ s$) Alfvén waves could be a viable candidate for the heating of quiet region loops. Further, Alfvén waves may be regarded as responsible for accelerating the solar wind.

Hollweg (1982a, 1985b) studied a train of switch-on shocks formed in the chromosphere from a train of Alfvén waves. Such trains of shocks give rise to a heating rate ($erg\ cm^{-3}\ s^{-1}$)

$$\epsilon_{SW} = \frac{B_0^2}{32\pi P}\left(\frac{\Delta v}{v_A}\right)^4 \quad , \tag{4.15}$$

where Δv is the transverse velocity jump at the shock and P is the wave period. The effect scales as B_0^{-2} and is effective only in weak field regions, like coronal holes. If $P = 100\ s,\ B_0 = 5\ G,\ \rho = 3.3 \cdot 10^{-16}\ g/cm^3$ and $\Delta v = 200\ km/s$, Hollweg finds $\epsilon_{SW} = 1.0 \cdot 10^{-5}\ erg\ cm^{-3}\ s^{-1}$. However, as the observed velocities Δv are about ten times smaller (Withbroe 1988), ϵ_{SW} should be decreased by a factor $\approx 10^4$.

Hollweg et al. (1982) study the nonlinear time-dependent adiabatic propagation of initially purely torsional Alfvén waves in open vertical solar magnetic flux tubes. They find that these waves can steepen into switch-on shocks in the chromosphere and can pass through the transition region into the corona to heat it. In addition mode-coupling to longitudinal waves is found. The longitudinal wave shows a slow shock as well as the fast shock which coincides with the switch-on shock. They suggest that the Alfvén waves drive upward flows which may explain spicules and may heat the upper chromosphere.

A different mechanism of damping of Alfvén waves has been investigated by Lou and Rosner (1986). According to them, when the Alfvén waves propagate in an infinitely conducting, inviscid, cold, incompressible and turbulent plasma in the presence of a strong magnetic field the waves get damped via scattering due to the background fluctuations. They find that the long period Alfvén waves may be effectively damped (thermalized) on time scales of the order of the (longitudinal) Alfvén crossing time if the background fluid is weakly turbulent as is the case with solar coronal loops.

Zähringer and Ulmschneider (1987) as well as Ulmschneider and Zähringer (1987) compute the propagation of nonlinear time- dependent adiabatic longitudinal and transverse mhd waves along a thin vertical flux tube in the solar atmosphere. The waves are excited at the foot of the tube by pure transverse shaking. Strong mode-coupling to longitudinal waves is found. The tube mass is seen to be lifted leading to adiabatic cooling of the tube. The lifting of the mass is attributed to centrifugal forces generated by the swaying of the tube. The authors conclude that mode-coupling appears to be important for chromospheric and coronal heating (see also Ulmschneider 1986).

Similon and Sudan (1989) study the energy dissipation of Alfvén wave packets in coronal arches. The authors find that more complex stochastic magnetic geometries which are expected to be encountered in realistic situations will dramatically decrease the dissipation length.

Wentzel (1989) studies the conversion of Alfvén waves to fast- mode waves and their subsequent Landau damping near the height where the coronal-hole nozzle diverges rapidly. Refraction can cause the conversion of up to half the Alfvén wave energy to fast-mode energy.

4.4. Dissipation via Landau Damping

Stéfant (1970) studied collisionless damping of shear Alfvén waves in the limit of low frequency and small but finite Larmor radius. Because the averaging of the wave electric field over the Larmor circle (the finite cyclotron radius effect) creates a small difference in the transverse velocities of ions and electrons, a longitudinal electric field appears in the Alfvén wave. This electric field is responsible for the transfer of wave energy into the thermal motions of the resonant particles. He calculates a normalized damping rate for shear Alfvén waves and for ion-acoustic waves and finds that the damping for the Alfvén waves is a maximum when the Alfvén velocity is equal to the real part of the phase velocity of the ion-acoustic wave.

Hollweg (1971) shows that large amplitude linearly or elliptically polarized Alfvén waves propagating parallel to the magnetic field can be dissipated by non-linear Landau damping which is caused by the longitudinal electric field associated with the ion-sound wave driven in second order by the Alfvén wave.

For additional work on Landau damping see Sections 3.4 and 3.5.

4.5. Dissipation by phase-mixing

Heyvaerts and Priest (1983) investigate coronal heating by phase-mixing both for open fields and for loop situations. The authors consider wave propagation in the z-direction of a cartesian coordinate system assuming an equilibrium magnetic field distribution $B(x)\hat{e}_z$ which varies in x-direction. $z = 0$ is the photospheric boundary. In the closed loop situation an additional boundary is assumed at $z = l_{\parallel}$. The fluid motions are in y-direction to simulate torsional Alfvén waves. Coherent harmonic excitation occurs at the foot points by prescribing the velocities $v_1(x,t)\hat{e}_y$ at $z = 0$ and $v_2(x,t)\hat{e}_y$ at $z = l_{\parallel}$. $\hat{e}_y$ and $\hat{e}_z$ are the unit vectors in y and z-directions. Assuming the coronal fluid to be weakly dissipative the wave equation is given by

$$\frac{\partial^2 v}{\partial t^2} = v_A^2(x)\frac{\partial^2 v}{\partial z^2} + \nu\left(\frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial z^2}\right)\frac{\partial v}{\partial t} \quad , \tag{4.16}$$

where ν is the sum of kinematic viscosity and ohmic diffusivity. In open configurations, neglecting dissipation, the coherent excitation of fixed frequency will propagate with the velocity $v_A(x)$, and as the altitude z grows these oscillations will get more and more out of phase. In closed configurations, it is the parallel wavelength which is fixed, and stationary oscillations excited in this resonant cavity progressively phase -mix as time elapses. These oscillations are coupled by weak friction.

In order to solve Eq. (4.16) for a coronal loop one writes

$$v(x,z,t) = v_1(x,t) + \frac{z}{l_{\parallel}}\left[v_2(x,t) - v_1(x,t)\right] + w(x,z,t) \quad . \tag{4.17}$$

where w is an arbitrary velocity function which is zero at both $z = 0$ and $z = l_{\|}$. w is written in a Fourier series in z on the interval $(0, l_{\|})$

$$w(x,z,t) = \sum_{m=1}^{\infty} b_m(x,t) \ \sin \frac{2m\pi}{l_{\|}}(z - l_{\|}/2) \quad . \tag{4.18}$$

The time-dependent Fourier coefficients $b_m(x,t)$ of this expansion must satisfy the relation

$$\left(\frac{\partial^2}{\partial t^2} + \Omega_m^2(x) \ - \nu \frac{\partial^2}{\partial x^2} \frac{\partial}{\partial t} \right) b_m = \frac{(-1)^{m+1}}{m\pi} \frac{\partial^2}{\partial t^2} \left[v_1(x,t) - v_2(x,t) \right] \quad , \tag{4.19}$$

where

$$\Omega_m = \frac{m \ 2\pi v_A(x)}{l_{\|}} \quad . \tag{4.20}$$

Eq. (4.19) represents a continuum of forced oscillators weakly coupled by friction. In the limit of small damping and large phase-mixing ($\Omega_m(x) \ t >> 1$) the approximate solution is given by

$$b_m(x,t) = \int_{-\infty}^{t} dt' \ f(x,t') \ \frac{\sin \Omega_m(x)(t-t')}{\Omega_m(x)} \exp\left[-\frac{\nu}{6} \ (\frac{d\Omega_m}{dx})^2 \ (t-t')^3 \right] \quad , \tag{4.21}$$

where $f(x,t)$ is the RHS of Eq. (4.19). Thus the waves of frequency $\omega = \Omega_m(x)$ damp in a time

$$t_D = 6^{1/3} \ \nu^{-1/3} \ (\frac{d\Omega_m}{dx})^{-2/3} \quad , \tag{4.22}$$

which may be rewritten as

$$t_D = \left(\frac{6 \ \Delta\omega \ \Delta a^2}{\nu} \right)^{1/3} \frac{1}{\Delta\omega} \quad . \tag{4.23}$$

In obtaining Eq. (4.23) use has been made of $d\Omega_m/dx = \Delta\omega/\Delta a$, where Δa is the thickness of the inhomogeneous region. $(\Delta\omega)^{-1}$ is the phase coherence loss time. The dimensionless quantity in the parenthesis is the Reynolds number $R_e \approx 6\Delta\omega\Delta a^2/\nu$ and has a typical value $\simeq 10^5$ for coronal loops. According to Heyvaerts (1985) the damping time comes out to be about 10 periods and represents an upper limit.

We find Heyvaerts estimate overly pessimistic. On basis of the values by Golub et al. (1980), see Sec. 10, for active region loops we assume, following Hollweg (1985b), an Alfvén speed variation from $v_{A1} = 3000 \ km$ to $v_{A2} = 2000 \ km$ over $\Delta a = 1000 \ km$. With $\nu = 10^{-16} \ T^{5/2}/\rho = 2 \cdot 10^{14} \ cm^2/s$, where we used $\rho = 5 \cdot 10^{-15} \ g/cm^3$ and $T = 2.5 \cdot 10^6 \ K$ we obtain from Eq.(4.22), $t_D = 67 \ s$. The resonant period is given by Eq. (4.6) as $P_{res} = 160 \ s$. Here we used the average length $l_{\|} = 2 \cdot 10^5 \ km$ for the active region loops of Golub et al. (1980). For a wave flux of $F_M = 1 \cdot 10^7 \ erg \ cm^{-2} \ s^{-1}$ by Withbroe and Noyes (1977) for active regions we get a heating rate $\epsilon_H = F_M/(v_A t_D) \approx 6 \cdot 10^{-4}$ *erg* cm^{-3} s^{-1} which is quite adequate for active region loops (see Table I as well as Hollweg and Sterling 1984). This shows that phase-mixing is a viable mechanism.

Heyvaerts and Priest (1983) derive an average energy dissipation rate ($erg\ cm^{-3}\ s^{-1}$) given by

$$\epsilon_H = \frac{1}{2}\,\rho(x)\,\nu \int_0^{l_\parallel} \frac{dz}{l_\parallel}\left(\frac{\partial v}{\partial x}\right)^2 = \sum_{m=1}^{\infty} \frac{1}{4}\,\rho(x)\,\nu\left(\frac{\partial b_m}{\partial x}\right)^2 \quad . \qquad (4.24)$$

Since the photospheric motions represent a stationary random process they will have some velocity power spectrum, $S_p(x,\omega)$. ϵ_H can be expressed in terms of these spectral contributions. If the width of the coronal loop in the y-direction is $l_\perp$ then the total heating rate (erg/s) in the loop structure is eventually found by integrating over the horizontal extent of the heating layer:

$$\epsilon_H\ l_\parallel\ l_\perp^2 = \frac{\mu}{\pi}\ l_\parallel\ l_\perp \sum_{m=1}^{\infty} \int_{-\infty}^{\infty} dx\ \rho(x)\ \Omega_m^2(x) S_p\big(x, \Omega_m(x)\big) \quad . \qquad (4.25)$$

Here μ is a pure number of order unity. Eq. (4.25) reveals that in the limit of strong phase- mixing and weak damping the total heating rate does not depend on viscosity (and resistivity), very similar to the resonant heating picture of Ionson (1982). The heating rate depends only on the properties of the photospheric driving source. Resonant heating is also discussed in Sec. 7.

According to Browning and Priest (1984) a shear Alfvén wave propagating in a laterally inhomogeneous structure develops strong velocity gradients due to phase-mixing. The strong gradients are subject to Ohmic and viscous dissipation so that phase-mixing may greatly enhance the damping of Alfvén waves and thus provide a viable mechanism for coronal heating.

Nocera et al. (1984, 1985) as well as Nocera and Priest (1984) study the phase-mixing of propagating Alfvén waves in detail and develop further the simplified treatment of Heyvaerts and Priest (1983).

Sakurai and Granik (1984) study the resonance of Alfvén waves in a coronal loop excited at its footpoints by the photospheric convective motion. They take into account the non-uniformity of the Alfvén velocity in the loop so that the Alfvénic surface wave resonance and the phase-mixing of Alfvén body waves may take place. Their wave energy dissipation rate turns out to be sufficient to heat the corona. It neither depends on the values of resistivity and viscosity nor on whether the Alfvén velocity is uniform or not. The velocity amplitude in the loop, however, depends on resistivity and viscosity as well as on the degree of non-uniformity.

Abdelatif (1987) investigates the dissipation by phase-mixing of shear Alfvén waves in a coronal loop driven externally by an incident wave in the subcoronal region. He finds that the total energy deposited in the loop depends on the magnetic diffusivity and viscosity, contrary to the conclusion of Ionson (1982) and Heyvaerts and Priest (1983). In addition he criticizes the boundary conditions applied at the loop foot points in the resonant heating picture by Ionson (1982). Abdelatif computes the energy deposited in a three-layer model by incident waves with periods of 5 minutes and 5 seconds and finds that 5 minute waves deposit almost no energy in small loops but 5 second waves deposit a substantial amount of energy in the loop although not sufficient to account for the heating of small loops.

Poedts and Goossens (1987, 1988) investigate poloidal wavenumber coupling of ideal mhd continuum modes in two-dimensional models for coronal loops and arcades. They consider two physical causes for this coupling: cross-section variation of the loops/arcades and variation of the equilibrium density across and along the magnetic field lines. According to them, for realistic values of the parameters the continuous spectrum is modified such that the ranges of the continuum frequencies are considerably enlarged, and the derivatives of the continuum frequencies normal to the magnetic surfaces are considerably increased. Consequently, the phase-mixing time is reduced, and the efficiency of phase-mixing as a heating mechanism for solar loops and arcades is increased. They find that the dissipation of wave energy is larger at the top of the coronal loops, in agreement with observations.

Phase-mixing and some cases of mode-coupling as well as the process of resonant absorption discussed in Sec. 6 are examples of heating mechanisms which depend on an Alfvén speed variation transverse to the field. Mode-coupling and phase- mixing can arise in similar situations. In the case were the magnetic field is in z-direction, and the Alfvén speed changes in x-direction, phase-mixing occurs when the coherent B-field perturbations are in y-direction, while mode-coupling happens when the B-field variations are in x-direction. The latter motions will lead to compressions in slabs of width Δx and thus mode-couple to fast and slow-mode waves. Therefore, in regions where phase-mixing is important, mode-coupling should be important too. We conclude that, as discussed above, a changing or geometrically complex field geometry greatly helps to dissipate Alfvén waves.

To summarize this section, it appears that for the chromosphere Alfvén wave heating by nonlinear mode-coupling or by turbulent heating is important. For coronal loops resonant heating, phase- mixing, turbulent heating and mode-coupling are the promising heating mechanisms while for coronal holes phase mixing, turbulent heating, mode-coupling and Landau damping seem to be important.

5. Heating by surface Alfvén waves

Surface waves can exist whenever there are boundaries in a medium. The surface waves on a lake are a well-known phenomenon. Consider the hot plasma containing solar coronal loops where $\beta = 8\pi p/B^2 < 1$. The essential feature of these loops is that because of a difference in the matter density and magnetic field strength of adjacent loops the Alfvén speeds inside and outside the loops are different. In response to large scale shaking of the foot of the loop by turbulent photospheric motions in addition to a continuum of body Alfvén waves, Alfvénic surface waves are excited which travel along the loop. Their amplitude decreases exponentially inwards and outwards from the loop boundary. When the surface waves are undamped the energy flux is parallel to the magnetic field B of the loop. If one considers heating of coronal loops by surface waves it must be kept in mind that since the surface waves occupy only the boundary regions of the loop the energy flux requirement is more severe for this wave type than for body waves which have the entire loop cross section available for the energy propagation. Note, however, that the energy transport requirements for surface wave heating is less severe than that for current heating (c.f. Sec. 6) which is confined to extremely thin current sheets, sheaths or filaments.

The phase velocity v_{AS} of Alfvén surface waves (see e.g. Ionson 1978, Wentzel 1978, 1979a,b,c) is intermediate between the Alfvén speeds on the two sides of the loop boundary:

$$v_{AS} = \frac{\omega}{k_{\|}} = \left(\frac{B_i^2 + B_o^2}{4\pi(\rho_i + \rho_o)} \right)^{1/2} \quad , \tag{5.1}$$

valid if $k_{\|}/k_{\perp} << 1$ and $\beta < 1$ and where i, o refer to the regions inside and outside the loop. $\|$ and $\perp$ indicate directions parallel and perpendicular to the loop. For example, with $v_{Ai} = 3000\ km/s$, $v_{Ao} = 2000\ km/s$, where v_{Ai} and v_{Ao} are the Alfvén speeds inside and outside the loop, and for simplicity $\rho_i \approx \rho_o$, one finds $v_{AS} = 2550\ km/s$. For wave periods $P = 2\pi/\omega = 300, 100, 10\ s$ one gets wavelengths $\lambda = 2\pi/k_{\|} = 7.7 \cdot 10^5, 2.6 \cdot 10^5, 2.6 \cdot 10^4\ km$, respectively. This shows that only for $P \leq 20\ s$ is λ less than the length $l_{\|} = 10^5\ km$ of a large loop.

If the boundary of the loop is very sharp the Alfvénic surface waves travel along it with little dissipation. Usually the boundary is diffuse and the boundary layer in which the Alfvén speed is changing is thin compared to the loops transverse diameter $l_{\perp}$. In this case after Ionson (1978) the surface wave resonates with a body Alfvén wave, called kinetic Alfvén wave, at a site (in the boundary layer) where the phase speed of the surface wave matches the local Alfvén speed. The kinetic Alfvén wave receives energy from the surface wave and its amplitude grows considerably.

This phenomenon is called *resonant absorption.* Note that resonant absorption refers to a spatial resonance which does not depend on the loop length and thus should not be confused with the temporal resonance discussed in Sec. 4. The latter resonance arises from constructive interference effects and lives from the reflections at the end points of a loop of given length, leading to standing waves (c.f. Eq. 4.6). The kinetic Alfvén waves dissipate in a thin sheath, the so called resonant absorption layer around the loop via Joule- and viscous heating. Thus surface Alfvén waves eventually heat the loop through Joule- and viscous dissipation.

The idea of spatial resonance and plasma heating was put forward by Grossmann and Tataronis (1973) as well as Hasegawa and Chen (1974). It was developed further by Chen and Hasegawa (1974) as well as Hasegawa and Chen (1976). Ionson (1978) used this idea to explain heating of "coronal rain" or sunspot loops. The damping rate (s^{-1}) of the Alfvénic surface wave is found to be (Lee and Roberts 1986, Hollweg 1985b, 1987a, 1987c)

$$t_{AS}^{-1} = \frac{\pi\ \Delta a\ k_{\|}^3\ |v_{Ai}^2 - v_{Ao}^2|}{8\ \omega} \quad . \tag{5.2}$$

Here $2\Delta a$ is the thickness of the boundary layer in which the Alfvén speed changes. For the above assumptions of the Alfvén speed and using $\Delta a = 1000\ km$ one finds $t_{AS} = 1.9 \cdot 10^4, 2.2 \cdot 10^3, 22\ s$ for $P = 300, 100, 10\ s$, respectively.

The thickness of the heated sheath is roughly given by (Ionson 1978)

$$\Delta x_s \approx R_{gi}\left(\frac{\Delta a}{R_{gi}}\right)^{1/3} \quad , \tag{5.3}$$

where $R_{gi} \approx 1.51\, T_{sh}^{1/2} B^{-1}$ is the ion gyroradius. For a long loop with $\Delta a = 10^4 km$ and $B = 10\ G$ one obtains $\Delta x_s < 1\ km$. A short loop with $\Delta a = 10^3 km$ and $B = 100\ G$ would have a sheath thickness of 0.1 km. T_{sh} is the sheath temperature and B is the loop's magnetic field strength.

According to Ionson (1978) the intense heating in the thin sheath at the boundary of the loop will drive upward convection in a layer next to the sheath whose thickness is given by

$$\Delta x_c \approx 3.5 \cdot 10^{-5}\, T_{sh}^{3/8}\, l_{\|}^{3/8}\, B^{-1/2} \quad , \tag{5.4}$$

where $l_{\|}$ is the loop length. The plasma rises in this convective layer and being continuously replaced extracts the heat from the sheath (see also Hollweg 1981a). The thickness Δx_c of the convection layer turns out to be about 30 km for long coronal loops and about 3 km for the short loops. The upward flow speed in the convection layer is roughly the free-fall speed which is about 100 km/s for a long loop. When the upflowing hot convection layer plasma reaches the top of the loop, it is transferred to the interior of the loop due to a Rayleigh-Taylor instability, thereby losing contact with the hot sheath. It begins to cool via radiation near the top and subsequently falls in a cool core whose diameter is essentially the loop diameter. Ionson offers this downflow as an explanation for "coronal rain" loops.

The temperature, T_{int}, of the cool core is given by Ionson (1978) as

$$T_{int} \approx 2.4 \cdot 10^{12} \frac{T_{sh}}{n_e\, l_{\|}^{1/2}} \quad , \tag{5.5}$$

where n_e is the electron density. For $T_{sh} = 3 \cdot 10^6 K, n_e = 7 \cdot 10^9 cm^{-3}, l_{\|} = 1 \cdot 10^{10} cm$, one obtains $T_{int} \approx 10^4 K$.

The surface Alfvén wave heating rate ($erg\ cm^{-3} s^{-1}$) is given by (Hollweg 1985b)

$$\epsilon_{AS} \;=\; 2\,\rho\,\delta v_{rms}^2 / t_{AS} \quad . \tag{5.6}$$

Ionson (1978) found that this rate is independent of the detailed viscosity or resistivity coefficient. Taking $v_{rms} = 20\ km/s$ and $\rho = 5 \cdot 10^{-15}\ g/cm^3$ one obtains $\epsilon_{AS} = 2.1 \cdot 10^{-6}, 1.8 \cdot 10^{-5}, 1.8 \cdot 10^{-3}\ erg\ cm^{-3} s^{-1}$, for waves of period $P = 300, 100, 10\ s$, respectively. This heating occurs in the loop boundary regions where one has a surface wave flux of $\rho \delta v_{rms}^2 v_{AS} = 5.1 \cdot 10^6\ erg\ cm^{-2} s^{-1}$.

It may be remarked that because the entire energy of the surface wave must be thermalized in a very thin dissipation sheath the excursions of the material near the sheath is much larger than the sheath thickness and hence nonlinear effects such as reconnection, shock formation etc. may lead to an effective nonlinear broadening of the dissipation sheath (Hollweg 1981a, Wentzel 1981). In addition the above picture applies only for (sunspot-) loops with cool cores. Generally coronal loops have hot cores (see Table I). For hot loops the above picture does not apply, as it would be difficult to explain how the interior should attain a higher temperature than the sheath heated by the surface wave.

There has been some discussion about the interpretation of the "surface wave damping rate". Ionson (1978) and Wentzel (1979b) start with a pure mhd system which is strictly dissipationless but they obtain a non-zero value of the damping rate. Lee (1980) as well as Lee and Roberts (1986) suggest that the wave decay rate should be interpreted as redistribution of wave energy in space or a mode-conversion rate. It appears now that the mode conversion rate can be interpreted as a plasma heating rate, at least in presence of some dissipative agents (Heyvaerts 1985, Hollweg 1987a,b,c).

It is not possible to observe the heated sheath (thickness $\leq 1\ km$) with the present resolution of earth-based or space-borne instruments, even the convective layer of 3 - 30 km thickness could not be observed. Ionson's (1978) predictions do agree with observed hot (sunspot-) loops with cool cores, at least qualitatively (Foukal 1976, Levine and Withbroe 1977). Alfvénic surface waves appear to supply the required energies to the coronal active regions with coronal velocities of the same order (tens of km/s) as are inferred from non-thermal line widths. To confirm the proposal of heating by Alfvénic surface waves the non-thermal emission resulting from the runaway electrons created by the parallel electric field of the resonantly excited kinetic Alfvén wave in the sheath should be investigated both theoretically and observationally.

Uberoi (1982) derives a relationship between the Alfvén speeds v_{Ai} and v_{A0} on the two sides of the boundary which is required for the existence of Alfvén surface waves in a low- β plasma. Somasundaram and Uberoi (1982) study compressibility effects on hydromagnetic surface waves. Roberts and Mangeney (1982) as well as Roberts (1983) discuss solitons in flux tubes, along with surface sausage waves.

Gordon and Hollweg (1983) investigate the collisional damping of surface Alfvén waves by viscosity and heat conduction. The authors consider surface waves at tangential discontinuities with a jump in Alfvén speed and do not take into account the effects of spatial resonant absorption proposed by Ionson (1978). They find that surface waves dissipate efficiently if their periods are shorter than a few tens of seconds and the background magnetic field is less than about 10 G.

Hasegawa (1985) presents mechanisms of Alfvén wave heating in space-astrophysical plasmas with particular emphasis to the parallel electric field generated in the mhd perturbations due to finite Larmor radius effects. Further discussion of this work can also be found in Hasegawa and Uberoi (1982).

Steinolfson et al. (1986) examine viscous damping of Alfvén surface waves both analytically and numerically using an incompressible mhd approximation. They show that normal modes exist on discontinuous as well as on continuously varying interfaces in Alfvén speed. These waves experience negligible decay below the chromosphere-corona transition layer. High frequency waves damp just above the transition layer, while those of lower frequency lose energy further out in the corona. By comparing the dissipative decay rates they find that wave damping by viscosity proceeds approximately two orders of magnitude faster than by resistivity. The importance of viscosity has also been stressed by Hollweg (1985b).

Assis and Busnardo-neto (1987) calculate the Cherenkov damping (Landau damping and transit time magnetic pumping) of waves with an arbitrary vector, in particular that of low frequency surface waves. According to them, a surface wave has a zeroth-order magnetic field parallel to the magnetic field of the loop and it is this component of the wave field which leads to transit time magnetic pumping in a collisionless plasma. In some situations surface waves can deposit energy and momentum in the corona via transit time magnetic pumping as efficiently as kinetic Alfvén waves do via Landau damping.

Davila (1987) investigates the heating of the solar corona by resonant absorption of surface Alfvén waves. He calculates the heating rate by an improved method, compares his results with observations and concludes that resonant absorption is a viable mechanism for the heating of the solar corona.

Mok (1987) as well as Einaudi and Mok (1987) study dissipation by viscous and resistive damping of surface Alfvén waves in a non-uniform plasma by using a normal mode analysis following the procedure of Mok and Einaudi (1985). Dissipation is important only in the narrow resonant layer where the eigenfrequency matches the local Alfvén frequency. The effects of viscosity are found to dominate resistivity in the quiet-sun corona and the solar wind. They obtain an energy deposition rate which is independent of the detailed dissipation processes via Joule or viscous dissipation, in agreement with earlier findings (Ionson 1978, 1982, Heyvaerts and Priest 1983, Hollweg 1984b, Lee and Roberts 1986). The authors conclude that dissipation by resonant absorption is a viable coronal heating mechanism.

Hollweg (1985b, 1987a,b,c) studies resonant absorption of surface waves using the incompressible mhd approximation both in simple terms, employing the thin flux tube approximation, and in detail. He shows that the energy of the surface wave propagating at a boundary layer of thickness $2\Delta a$ gets deposited into a much thinner "energy containing layer" of thickness $k_{\|}(\Delta a)^2$ near the resonant field line within the boundary layer. This confirms similar findings of Lee and Roberts (1986). In an investigation which explicitly includes the presence of viscosity the net steady-state heating is found to be independent of viscosity, in agreement with Ionson (1978). He points out that if his incompressible case results could be extrapolated to the compressible case then resonant absorption must be considered as a viable mechanism for coronal heating.

Grossmann and Smith (1988) study resonant absorption of a spectrum of standing Alfvén waves in coronal loops. The input spectrum has a temporal part consisting of granular and 300 sec type solar p-mode oscillations and a spatial part for which a power law was assumed. For cylindrical tubes of various radial density contrasts (between the tube axis and the boundary), of various helicities and of various radial Alfvén crossing times it is shown which band of the input spectrum is absorbed. The heating is confined to the boundary region of the loop. The authors conclude that resonant absorption of Alfvén waves is a viable mechanism for coronal loop heating.

To summarize, resonant absorption of surface Alfvén waves appears to be a viable heating process both for open regions and for coronal loops.

6. Heating by current (or magnetic field) dissipation

Observed coronal loops can often be fitted quite well to potential fields with $\nabla\times\mathbf{B} = 0$ and the potential continuation of the longitudinal component of the photospheric magnetic field distribution over the solar surface reproduces quite well the gross properties of the large scale solar topology. Even where the coronal topology appears potential, however, small departures from the potential form would be difficult to detect since individual loops are often diffuse and only barely resolved. In other cases the topology is more consistent with force-free fields described by

$$\nabla \times \mathbf{B} = \alpha \mathbf{B} \quad , \tag{6.1}$$

where α is a constant. As force-free fields and departures from potential fields lead to currents, it is thus widely believed that electric currents must exist in the corona and that these currents are an important source of coronal heating. However, as heating by volume currents is much too small, heating in current sheets must be considered.

6.1. Heating flux

Foot point motions due to the granular and supergranular flows as well as due to the differential rotation can lead to a buildup of the magnetic energy content of coronal loops. Consider a coronal loop of length $l_{||}$ (foot to foot) with the magnetic field B_0 along the loop. Radial and torsional photospheric foot point motions with velocity v will displace the field lines by a distance $l_\perp \approx vt$ in a characteristic time t. This generates a magnetic field perturbation $B_1 = B_0 l_\perp / l_{||}$ and consequently an energy density (erg/cm^3)

$$B_1^2/4\pi \;\approx\; \frac{B_0^2}{4\pi}\left(\frac{v}{l_{||}}\right)^2 t^2 \quad , \tag{6.2}$$

For slow foot point motion, where $t > t_A \equiv l_{||}/v_A$, the heating flux $(erg\ cm^{-2}\ s^{-1})$ entering the loop is given by

$$F_{H,DC} = \frac{B_1^2}{4\pi}\frac{l_{||}}{t} \approx \frac{B_0^2}{4\pi}\frac{v^2}{l_{||}}\, t \quad , \tag{6.3}$$

while for fast foot point motion, $t < t_A$, where the effective loop length is $l_{||} = v_A t$, the heating flux is given by

$$F_{H,AC} = \rho v^2\, v_A \quad , \tag{6.4}$$

where in the above equations v_A is the Alfvén speed and t_A the Alfvén transit time of the loop (Leibacher and Stein 1982, Parker 1983b, 1986). Here DC refers to the direct currents involved in the energy dissipation in coronal loops as discussed in this Section. AC refers to the alternating currents associated with Alfvén waves. Heating by these Alfvén body and surface waves has already been discussed in Sections 4 and 5.

The stochastic processes which lead to the energy input into the corona after Eq. (6.3) have been investigated in greater detail by Sturrock and Uchida (1981). They consider fluctuating twisting velocities $v_\phi = Rd\chi/dt$, where R is the loop radius and χ the twisting angle. The energy flux entering the corona is given by

$$F_H = \frac{B_0^2 R^2}{16\pi l_{\|}} \left\langle \frac{(\Delta\chi)^2}{\Delta t} \right\rangle \quad , \tag{6.5}$$

where the expected twist build-up can be expressed as

$$\left\langle \frac{(\Delta\chi)^2}{\Delta t} \right\rangle = \frac{4}{R^2} < v_\phi^2 > t_c \quad . \tag{6.6}$$

Here t_c is a correlation time and the twist from both feet of the loop is considered.

6.2. Simple Joule heating

Gold (1964) pointed out that in the turbulent $\beta \geq 1$ plasma below the photosphere, the magnetic field has a large curl as compared to the small or zero curl of the coronal field. He argued that upon entering the upper atmosphere the subphotospheric fields must "shed their curls" by dissipating the associated current in the chromosphere and inner corona.

Two mechanisms capable of dissipating currents are discussed in the literature, *simple Joule heating* and *magnetic reconnection*. In simple Joule heating no topological changes in the magnetic flux surfaces of the structure take place. In reconnection there are topological changes associated with strong convective flows and Joule heating concentrated in current sheets.

A force-free non-potential magnetic field **B** can be associated with a current density $\mathbf{J}_{\|}$ parallel to the magnetic field which is given by

$$\mathbf{J}_{\|} = -n_e\, e\, v_d = \frac{c}{4\pi} \nabla \times \mathbf{B} = \frac{c\alpha}{4\pi} \mathbf{B} \quad , \tag{6.7}$$

where n_e is the electron density, e the electronic charge and v_d the drift velocity of the electrons. Due to the resistivity η of the medium the current produces a heating rate ($erg\ cm^{-3} s^{-1}$)

$$\epsilon_H = \eta J_{\|}^2 \quad . \tag{6.8}$$

Kopecky and Obridko (1968) note that the classical resistivity, due to Coulomb collisions between electrons and ions, is too small to provide significant heating of the corona. The only way the magnetic field or current dissipation can produce the necessary coronal heating is for the magnetic field changes and accompanying electric currents to be concentrated in extremely intense current sheets, current sheaths or current filaments (see e.g. Priest 1982). If the current density is large and the width of such a current distribution is less than a few meters, the dissipation may be considerably enhanced by the presence of plasma turbulence.

Tucker (1973), following the suggestion of Gold (1964) proposed that solar active regions are steadily heated by magnetic energy dissipation. According to him the magnetic energy is stored at the total rate (erg/s)

$$U_M \simeq \frac{v_\phi B^2}{8\pi} A \quad , \tag{6.9}$$

where v_ϕ is the velocity of twisting of the magnetic field of strength B by photospheric motions over an area A and at the same time it is being dissipated at the total rate

$$U_{diss} \simeq \eta J^2 \, Ah \quad , \tag{6.10}$$

for currents J distributed uniformly throughout the volume Ah of the active region. In order to have steady state heating both rates must be equal. Tucker finds that a uniform current distribution leads to a very high value of $\nabla \times \mathbf{B}$ ($\sim 10^{-2.4} G/cm$) which is inconsistent with observations. He thus concludes that dissipation must take place in thin current sheets. Assuming a single current sheet in an active region Kumar and Narain (1985) found unacceptably high values of the twisting velocity ($\sim 10^{14}\ cm/s$). In order to get acceptable values of the twisting velocity the number of current sheets in the active region should be large (Kumar 1987).

Somov and Syrovatskii (1977) also argue in favour of the proposal that the heating source for an active region is the quasi-steady dissipation of magnetic fields in current sheets. Syrovatskii (1978) as well as Bobrova and Syrovatskii (1979) show that for force-free fields complex foot point motions create current singularities which get rapidly dissipated.

Rosner et al. (1978a) extend Tucker's work and consider a coronal loop of length $l_\parallel$ and radius R heated in a thin sheath of thickness ΔR around the loop. From Eq. (6.7) they obtain

$$J_\parallel \simeq \frac{c}{4\pi} \frac{\Delta B}{\Delta R} \quad , \tag{6.11}$$

where ΔB is the magnetic field jump across the sheath. This allows to derive a heating rate ($erg\ cm^{-3}\ s^{-1}$) averaged over the entire loop

$$\epsilon_H = \eta J^2 \, \frac{2\pi R \, \Delta R \, l_\parallel}{\pi \, R^2 \, l_\parallel} \simeq \eta \frac{c^2}{8\pi^2} \, \frac{(\Delta B)^2}{R \, \Delta R} \quad . \tag{6.12}$$

This heating rate originating in the sheath is supposed to balance the average radiative cooling rate from the entire loop which in the thin plasma approximation can be written

$$\epsilon_R \simeq 7 \cdot 10^{-17} \, T^{-1} \, n^2 \quad , \tag{6.13}$$

where T is the temperature and n is the total number density. Assuming $\epsilon_H = \epsilon_R$ allows to solve for the sheath thickness

$$\Delta R \simeq 1.6 \cdot 10^{35} \eta \, n^{-2} T R^{-1} (\Delta B)^2 \quad , \tag{6.14}$$

which for reasonable values $R = 10^4\ km$, $\Delta B = 10\ G << B$, $T = 2.5 \cdot 10^6\ K$, $n = 2 \cdot 10^9\ cm^{-3}$ and a classical resistivity of $\eta_{class} = 10^{-7}T^{-3/2} = 3 \cdot 10^{-17}\ s$, leads to a value $\Delta R = 0.3\ cm$, which is of the order of the gyroradius of thermal electrons in a field of $B = 100\ G$ and thus is unacceptably small. In addition this small ΔR would imply large $J_{\|}$, and after Eq. (6.7) drift velocities v_d, that surpass the ion sound speed at which point various micro-instabilities as e.g. ion-acoustic turbulence are excited. Rosner et al. argue that the onset of ion-acoustic turbulence leads to a much higher anomalous resistivity $\eta_{ia} = 2 \cdot 10^{-8} n_e^{-1/2} = 6 \cdot 10^{-13}\ s$, through which the heating rate can be greatly magnified and the thickness of the sheath is increased to an acceptable value $\Delta R = 6 \cdot 10^3\ cm$.

However, the continued existence of a thin sheath may not be possible once the ion-acoustic turbulence is established, because there is enhanced diffusion and the sheath broadens as the electrons drift along it. Therefore the condition (6.14) may not be satisfied and the heating may stop unless the condition is restored by some external mechanism. Another difficulty is that the thin sheath can be disrupted by the double tearing mode in a time scale of the order of a few seconds (Hollweg 1981a, and references therein). The viability of anomalous current dissipation as a heating mechanism has been critically discussed by Chiuderi (1981) and Kuperus et al. (1981).

Duijveman et al. (1981) describe fast plasma heating by current dissipation via anomalous and inertial resistivity. The anomalous character of the current dissipation is caused by the excitation of electrostatic ion-cyclotron or ion-acoustic waves, the inertial resistivity by geometrical effects.

Spicer (1983) argues that regions of high current density cannot exist within the solar atmosphere in a quasi-stationary state if they do not already exist at the photospheric boundary. He points out that theoretical treatments of coronal heating by electrodynamic processes must take proper account of the photospheric spatial distribution of forces that generate the currents (waves) and not just the power contained in those waves that results in coronal heating.

Narain and Kumar (1985), as well as Kumar (1987) study anomalous current dissipation in flux-tubes of variable cross section which relaxes the constant tube cross-section assumption of Rosner et al. (1978a).

Hinata (1980, 1981) studies direct current dissipation in coronal magnetic loops. According to him a current-carrying loop will have non-symmetric temperature and density distributions across the top of the loop. When the electron drift velocity v_d is near or above the critical value $v_c = 0.3(T_i/T_e)^{3/2}v_e$ (where T_i and T_e are the ion and electron temperatures and v_e the electron thermal speed) for the onset of electrostatic ion-cyclotron turbulence, the plasma will be heated to a temperature of up to a few million K provided the electron density is $\sim 10^8 cm^{-3}$ and the height of the loop is $\sim 10^9\ cm$. Hinata (1983) suggests that small ($< 10^7\ cm$) chromospheric structures can be heated adequately by rather weak currents.

Sakurai and Levine (1981) study generation of electric currents in a magnetized plasma overlying a dense convective layer, assuming the magnetic field perturbation to be small and to satisfy a force-free equation. According to them currents are produced by rotational motions on the boundary in the case of a uniform equilibrium field. In a simple two-dimensional bipolar configuration both irrotational and incompressible motions give rise to currents, and the current density has a peak at the magnetic neutral line. They derive scaling laws for the current density as well as for the stored magnetic energy and discuss the possibility of heating the solar corona through the dissipation of the generated currents.

Ferrari et al. (1982) investigate the stability of a magnetized low-density plasma to current-driven filamentation instabilities and apply their result to the surface layers of stars.

Rabin and Moore (1984) propose a model for sun's lower transition region based on heating by filamentary currents that flow along the magnetic field. The current filaments have thicknesses in the range of 1 cm to 1 km and the potential magnetic field must be greater than about 10 G. They find that in the lower transition region, an ensemble of filamentary currents that agree in sign across the horizontal scale of a photospheric granule can generate enough heat to match observations. Heat transport takes place primarily by conduction across rather than along field lines. They do not mention the source of these currents.

Bodo et al. (1985) study the stability of magnetic loops to current- driven filamentation instabilities. They find that the mechanism for current filamentation could heat initially cold plasma entrained by magnetic flux emerging from the solar convection zone into the ambient photosphere and above. This mechanism is very effective at the top of magnetic loops. Bodo et al. (1987) investigate current-driven mhd thermal instabilities in sheared fields. According to them thermal instabilities in the presence of magnetic field naturally lead to the formation of filamentary structures. Anisotropy of thermal conduction in the presence of magnetic field aids this formation. Further the energy balance in the lower transition region between Joule heating and losses due to radiation and conduction across the magnetic field requires individual current filaments to be of thickness larger than $\sim 1cm$ and smaller than $\sim 1km$.

6.3. Magnetic reconnection

Parker (1972) shows that when two or more flux tubes wrap around each other to form a rope with braided or knotted flux tubes, the magnetic topology is without equilibrium, no matter what fluid pressures are applied along the individual lines of force. The result is rapid dissipation and field-line merging, which quickly reduces the topology to the simple equilibrium form. This process, i.e. "topological dissipation" could be a significant mechanism for the heating of the solar corona.

Levine (1974a,b) proposes that the corona is interspersed with magnetic neutral surfaces which collapse and accelerate particles (mainly ions) to a few times their thermal velocity. These accelerated particles travel through the corona and lose their increased energy through Coulomb collisions. This approach can qualitatively explain the existence of regions of enhanced (active regions, bright points) or decreased (coronal holes) heating in the corona. It can also explain the maximum and minimum corona qualitatively (Narain and Kumar 1986, Kumar 1987).

Priest and Raadu (1975) as well as Tur and Priest (1976) showed that current sheets can be formed about neutral points if by photospheric motions or by the emergence of new magnetic flux oppositely directed magnetic fields are pressed together. The dissipation of these sheets produces coronal heating.

Withbroe and Noyes (1977) point out that for the type of heating proposed by Levine the neutral sheets would have to be distributed through the large scale coronal structures at a scale small compared to the resolution of space-borne instruments or of ground based observations in order to explain the apparent smoothness and lack of rapid time variability of these structures. Chiuderi (1981) states that the original concept of Levine's model, reconnection in neutral sheets, should be replaced by that of localized tearing modes in current sheaths that appear to be more easily realized. According to Kuperus et al. (1981) Levine's proposal is qualitative in nature and has little observational support.

Glencross (1975), while explaining the heating of coronal material at X-ray bright points, suggested that merging of twisted braids is likely to occur to some extent in all magnetic regions of the corona and could be a significant source of coronal heating.

A somewhat different mechanism from that of Levine (1974a,b) has been suggested by Parker (1975) who investigated the stability of a magnetic flux tube as it rises through the photosphere to produce an X-ray bright point. He showed that if the flux tube is twisted more than a modest amount, equivalent to a 45 degree pitch over only one scale height, then no hydrostatic equilibrium exists and the tube buckles, spirals and/or lengthens into a contorted form which causes rapid dissipation and reconnection of the magnetic field because oppositely directed field elements are pressed together. Such a behaviour could be the cause of bright point flares which are observed to exhibit many characteristics of active region flares. If this process occurs continuously it could also produce quasi-steady heating of bright points or other magnetic flux tubes. This proposal has been subsequently developed further (Parker 1979, 1981a,b, 1982, 1983a,b,c, 1985, 1986, Vainshtein and Parker 1986).

The basic idea is that the non-equilibrium of the magnetic field produces current sheets which constantly grow thinner and more concentrated as the fluid is squeezed out so there is rapid reconnection and dissipation no matter how small the electrical resistivity is. Further the high electron conduction velocities within the current sheet may produce plasma turbulence, anomalous resistivity and resistive tearing mode instabilities (Parker 1986 and references therein).

Using Eq. (6.3), Parker (1983c, 1986) estimates that for an active region with $v = 0.5\ km/s$, $B = 100\ G$, $l_{\parallel} = 10^5\ km$ the time required to accumulate a degree of wrapping of the field lines which corresponds to an energy flux of $\sim 10^7\ erg\ cm^{-2}s^{-1}$ turns out to be about 14 hours. The same state is reached in an ephemeral region ($l_{\parallel} = 10^4\ km$) in about 1.4 hours. Parker believes that his dissipation mechanism destroys the current sheets as fast as they are created and thus could be a viable mechanism for heating the solar corona.

Golub et al. (1980) consider heating of closed coronal regions via twisting of the foot points. Using the scaling law of Rosner et al. (1978b) they find a relation $p \sim < B >^{1.7}$ which agrees with the observational data reasonably well. Here p represents gas pressure and B stands for the longitudinal field strength.

Galeev et al. (1981) discuss the heating of inhomogeneous corona based upon local coronal magnetic field energy release and associated plasma heating. This work represents an extension and refinement of previous work of Rosner et al. (1978a). Special emphasis has been given to non-linear kinetic tearing-mode reconnection. Derived relations qualitatively explain the substantial brightness variability of active region loop structures but relatively little variability in the derived average coronal temperatures.

Sturrock and Uchida (1981) by taking into account radiative and conductive losses derive a scaling law between the loop's gas pressure and the magnetic field strength $p \sim B^{6/7} l_{\|}^{-1}$ which, however, does not agree with that of Golub et al. (1980).

Steinolfson and Van Hoven (1984) study magnetic reconnection in the presence of fast thermal instability. If the radiative instability is faster than the tearing instability substantial magnetic field reconnection occurs which is explained by the large temperature drop or resistivity rise. They call this process radiative tearing.

Blake and Sturrock (1985) suggest that visible spicules result from the eruption of a small amount of magnetic flux that has become detached from the photosphere by reconnection. The eruption of magnetic flux into the corona, which is already permeated by a magnetic field, will lead to current-sheets at the interface between the old and new fields. If the magnetic field in the corona contains currents, it necessarily contains free energy which may be released by reconnection leading to heating of the corona outside active regions.

Van Ballegooijen (1985, 1986) reconsiders the suggestion of Parker (1972) that magnetic field, in general, does not attain equilibrium and must develop current sheets in which the braiding patterns (produced by convective flows below the solar surface) are dissipated (topological dissipation). He finds that the coronal field can simply adjust to the slowly changing boundary conditions in the photosphere and that topological dissipation of the winding patterns does not take place (see also Antiochos 1987). Instead he proposes that required current densities can be produced via a cascade process in which "free" magnetic energy is transferred from large to small length scales in the corona as a result of the random motions imposed on the magnetic field lines by subphotospheric convective flows. He obtains the following heating rate ($erg\ cm^{-3}\ s^{-1}$) for this process

$$\epsilon_H \simeq 0.19 B^2 D / l_{\|}^2 \quad , \tag{6.15}$$

where $D = (\pi/2)^{1/2} v^2 t$ is the effective diffusion constant of the photospheric motions. Adopting the values $D = 250\ km^2/s, B = 100\ G$ and $l_{\|} = 10^5\ km$ for an active region van Ballegooijen finds a heating rate $\epsilon_H \approx 5 \cdot 10^{-5}\ erg\ cm^{-3}\ s^{-1}$. Comparison with the empirical rate $\epsilon_R = 2\ F_M / l_{\|} = 2 \cdot 10^{-3}\ erg\ cm^{-3}\ s^{-1}$, where $F_M = 1 \cdot 10^7\ erg\ cm^{-2}\ s^{-1}$ has been taken from Withbroe and Noyes (1977) shows that ϵ_H is too low by a factor of ~ 40. However, ϵ_H is only a factor of 2 below the mean radiative losses $\epsilon_R = 1 \cdot 10^{-4}\ erg\ cm^{-3}\ s^{-1}$ estimated from Eq. (6.13) using $T = 2.5 \cdot 10^6\ K$ and $n = 2 \cdot 10^9\ cm^{-3}$.

Strauss (1988) calculates the reconnection rate in a current sheet in the presence of tearing mode turbulence. He finds a large heating rate for a current sheet which may be many orders of magnitude larger than that in a Sweet-Parker current sheet. He points out that hyperresistivity produced by tearing-mode turbulence may explain both fast reconnection and the heating of the solar corona.

Strauss and Otani (1988) notice that when the twisting of the coronal magnetic fields by the photospheric motions exceeds a critical amount, kink-ballooning instabilities occur which lead to the formation of current sheets. Reconnection and ohmic heating in these sheets in the presence of resistivity unwinds the twist of the magnetic field and thereby heats the coronal plasma.

6.4. Complex magnetic reconnection, magnetic helicity, turbulent heating

Taylor (1974) suggested that reconnection in a magnetic field takes place to such an extent that the new field finds itself in the minimum magnetic energy state compatible with the conservation of magnetic helicity, K, defined by

$$K \equiv \int_V \mathbf{A} \cdot \mathbf{B} \, dV \quad , \tag{6.16}$$

where $\mathbf{A}$ is the vector potential of the magnetic field $\mathbf{B}$ and the integral is carried over the volume in which the reconnection takes place. In perfect mhd, the magnetic helicity is an invariant for each closed flux tube (Woltjer 1958) and it has the property that the minimum energy state is some general force-free field. This could also be true for a dissipative system (Heyvaerts 1985, and references therein).

Heyvaerts and Priest (1984) as well as Browning (1984) have adopted the above idea to study the heating of the coronal loops by direct currents. They assume that at each moment of its evolution, complex reconnection phenomena occur very rapidly in the magnetic structure and that the system is always in the lowest energy state compatible with magnetic helicity which evolves according to the equation

$$\frac{dK}{dt} = \int_{boundary} (\mathbf{A} \cdot \mathbf{v})(\mathbf{B} \cdot d\mathbf{S}) \quad , \tag{6.17}$$

where v is the fluid velocity at the boundary. The right hand side is a surface integral enclosing the volume of the loop. Care must be taken because for an open-ended flux-tube K is not a gauge-invariant quantity. They arrive at the following expression for the heating flux of a coronal loop

$$F_H = \frac{B^2}{4\pi} \frac{v^2}{l_{\|}} t_D G_f \quad , \tag{6.18}$$

where t_D is the stress relaxation time and G_f is a geometrical factor given by

$$G_f = \left(\frac{l_{\|}}{l_{\|} + l_v} \right)^2 \quad , \tag{6.19}$$

where l_v is the length scale of velocity cells. This factor might be important when large velocity cells at the boundary are present. Except for the geometrical factor this flux is identical to Eq. (6.3).

In addition the analysis of Heyvaerts and Priest (1984) discusses an interesting border between flares and coronal heating. Both phenomena result from magnetic stress relaxation. Whereas coronal heating takes place when boundary motions are slow and stresses never accumulate very much while flares occur in configurations which evolve so quickly that stresses are not able to relax at the same pace as they are built up and ultimately relax violently (Heyvaerts 1985).

Browning et al. (1986) as well as Browning and Priest (1986) study heating of closely packed flux tubes and coronal arcades using the Taylor-Heyvaerts hypothesis. They find a heating formula quite identical to that of Parker (1983b). In case of coronal arcades they notice that similar foot-point motions heat an arcade more efficiently if it is already sheared, such as in a rapidly evolving active region. Vekstein (1987) extends the theory of Browning et al. (1986) and obtains an expression for the coronal heating rate that is valid for any given photospheric velocity field.

Dahlburg et al. (1988) by fully three-dimensional numerical simulation, study the time-dependent relaxation of a coronal gas column permeated by a force-free magnetic field, which initially has been disturbed by broad-band velocity and magnetic field perturbations. The system is seen to evolve through a succession of force-free states. The authors find that concentrated vortex structures develop in the helical flow which are the primary sites of turbulent heating. These vortex structures are almost as important for the heating as the electric current sheets which develop in the gas column.

Coronal heating by selective decay of MHD turbulence is also discussed by Gomez and Ferro Fontan (1988). The authors estimate that the cascade due to nonlinear interaction of turbulent modes is a viable coronal heating process.

7. Unified resonant LCR circuit approach

The photosphere (with $\beta \geq 1$) and the corona (with $\beta < 1$) of the sun are electrodynamically coupled since extrapolations of the observed photospheric magnetic energy into the corona exceeds the in situ thermodynamic energy. Observations show that a large part of the corona consists of loop-like structures. Ionson (1982), and more clearly derived in his 1984 paper, represented a magnetic loop system by an equivalent resonant electric circuit. Coronal heating is seen as the analog of the damping in a "series LCR" circuit.

Using the linearized MHD equations, assuming viscosity, finite electrical conductivity, finite phase locking and magnetic stress leakage into the $\beta \geq 1$ regions, Ionson derived a second order differential equation for the bulk current in the loop, which he compares with a similar equation for an LCR electric circuit. From this comparison the quantities L, C and R can be derived. According to Ionson the inductance, $L = 4l_{\|}/(\pi c^2)$, is a measure of the ability to store magnetic energy, the capacitance $C = l_{\|}c^2/(4\pi v_A^2)$, is a measure of the ability of the magnetoplasma to store electric and kinetic energy and the resistance $R = R_{tot}$, is a measure of the ability to convert electrodynamic energy into heat. His physical system consists of a $\beta < 1$ magnetic loop and an underlying photospheric region of $\beta \geq 1$ in which velocity fields v_θ generated by the convection zone drive electrodynamic activity within the loop via the interconnecting magnetic field. The equivalent electric circuit representation is

$$L\frac{d^2 I}{dt^2} + R_{tot}\frac{dI}{dt} + \frac{I}{C} = \frac{dE(t)}{dt} \quad , \tag{7.1}$$

where I is the bulk current. As the chromospheric and the photospheric resistances are small, Ionson finds that the resistance R_{tot} results only from the coronal part of the loop and that R_{tot} can be expressed by

$$R_{tot} = R_{diss} + R_{phase} + R_{leak} = L\Big(\frac{1}{t_{diss}} + \frac{1}{t_{phase}} + \frac{1}{t_{leak}}\Big) \quad , \tag{7.2}$$

where t_{diss} is the coronal energy dissipation time, t_{phase} is the phase-mixing time and t_{leak} is the magnetic stress leakage time. $E(t) = (l_{\perp} v_{\theta} B/c)_p$ is the driving e.m.f. which is generated by the velocity fields v_{θ}. Index p indicates photospheric values. Through Eq. (7.1), Ionson in a global manner incorporates the generation, propagation and dissipation of electrodynamic energy, which amounts to coupling photosphere, chromosphere and corona electrodynamically.

The equivalent LCR circuit approach leads to a characteristic resonant frequency, ω_0, given by

$$\omega_0 = (LC)^{-1/2} = \pi v_A / l_{\parallel} = 2\pi / t_A \quad , \tag{7.3}$$

where v_A is the average Alfvén speed and t_A is the coronal loops Alfvén transit time. For resonance heating the important quantity is the quality Q_{tot} of the LCR circuit given by

$$Q_{tot} = \frac{\omega_0}{2\,\Delta\omega} = \frac{\omega_0\, L}{R_{tot}} \quad . \tag{7.4}$$

Here $2\Delta\omega$ is the full width at half maximum of the resonant peak of I^2. Of the spectrum of Alfvén waves entering at the foot points of the coronal loop the LCR circuit absorbs the energy contained in the band $2\Delta\omega$. Thus, for resonance heating to be important one would like large $\Delta\omega$, that is large R_{tot} or small Q_{tot}. Ionson finds that heating explicitly depends upon the $\beta \geq 1$ velocity field's spectral power. Magnetic loops of different lengths are heated at a rate that critically depends upon the amount of spectral power available at their resonant frequencies. He finds that the power drain on the source is apparently totally independent of the mode of dissipation. Further he finds that the dominant form of dissipation in solar coronal loops is via ion-viscosity.

Martens and Kuperus (1982) as well as Martens and Kuin (1983), using Ionson's resonant electrodynamic theory, explain catastrophic changes in the loop's X-ray visibility (loop evacuation/coronal rain and loop brightening).

Ionson (1983) investigates diffuse and isolated heating in magnetically confined X-ray plasmas of astrophysical origin by using his LCR circuit representation. He points out that tenuous X-ray emitting plasmas ($T > 10^6 K$) are found in the sun, early-, late-type stars, accretion-disks, galaxies, galactic jets, active galactic nuclei etc. He assumes that the site of mechanical activity, characterized by $\beta \geq 1$ zones of differentially rotating, convective velocity fields couples to and energetically maintains a spatially distinct, yet contiguous site of $\beta < 1$ X-ray activity. He obtains relations for diffuse heating (resulting from efficient transport of energy throughout the magnetic loop volume) and isolated heating (within the electrodynamic dissipation shell) corresponding to Joule dissipation, shear-, and compressional-viscous dissipation as a function of the plasma temperature T, thermal pressure p and stressing velocity $v_{\beta \geq 1}$. These relations do not depend upon the magnetic field strength, explicitly.

Ionson (1984, 1985a,b) presents a unified theory of both resonant (AC) and nonresonant (DC) electrodynamic heating. Here the heating flux, F_H, is given by

$$F_H = F_{max}\, f_\epsilon \quad , \tag{7.5}$$

where F_{max} is the maximum flux of energy that could enter the coronal portion of the loop and is given by

$$F_{max} = 16\left(\frac{B}{B_p}\right)\left(\frac{v_{A,p}}{v_A}\right)\left(\frac{1}{2}\rho v^2\right)_p v_{A,p} \quad , \tag{7.6}$$

where index p indicates photospheric $\beta \geq 1$ regions. f_ϵ is an electrodynamic coupling efficiency describing the fraction that enters and heats the contained plasma; it is given by

$$f_\epsilon = \frac{Q_{tot}}{Q_{diss}}\frac{t_C}{t_A}\begin{cases} \left[1 + \left(\frac{t_{peak}}{t_A} - \frac{t_A}{t_{peak}}\right)^2 \left(\frac{t_C}{t_{peak}}\right)^2\right]^{-1} , & Q_{tot} > 1; \\ \left[1 + \frac{t_C}{t_A Q_{tot}}\right]^{-1} , & Q_{tot} < 1, \end{cases} \tag{7.7}$$

where t_C is the correlation time of the $\beta \geq 1$ convection and t_{peak} is the period at which the convection spectrum peaks. Further Q_{tot} and Q_{diss} are, respectively, the total and dissipative qualities of the loop and are given by

$$Q_{tot} = \frac{2\pi}{t_A}\left(\frac{1}{t_{diss}} + \frac{1}{t_{phase}} + \frac{1}{t_{leak}}\right)^{-1} \quad , \tag{7.8}$$

and

$$Q_{diss} = 2\pi \frac{t_{diss}}{t_A} \quad . \tag{7.9}$$

Ionson classifies heating theories in many categories such as high-quality coupling, low-quality coupling, stochastic coupling, dissipation limited, phase-mixed limited, leakage limited, etc. According to him this classification scheme consolidates a variety of new heating mechanisms and illustrates that Alfvénic surface wave heating, stochastic magnetic pumping, resonant electrodynamic heating, and dynamical dissipation are actually special cases of a more general formalism.

He further states that in case of solar coronal loops the total quality is essentially equal to the dissipative quality (i.e. $Q_{tot} \simeq Q_{diss}$) and the coupling category changes as loops age and hence increase in length. Specifically, young active region loops are found to have $Q_{diss} < 1$ and are heated at a rate that depends upon the dissipation time, t_{diss}, while older active region and large scale loops have $Q_{diss} > 1$ and are heated at a rate that is quantitatively independent of the dissipation time. He points out that active region loops appear to be heated by electrodynamical coupling to $\beta \geq 1$ p-mode oscillations, while large scale loops are heated by coupling to the solar granulation. Loops with lengths larger than the solar radius resonate with supergranulation time scales ($\sim$ 20 hours) and are quite rare (Ionson 1985a).

Mullan (1984) studies the possibility of resonant electrodynamic coupling in the coronae of red dwarfs. According to him the convective time scale, t_C, is expected to decrease significantly as one goes to later spectral types. In the sun the value of t_C is many hundred seconds but values as small as a few tens of seconds appear to be characteristic of the convection zones in the coolest M dwarfs. Information how t_A varies along the main sequence is scarce but it appears that it increases towards cooler stars. Therefore a crossover of two scales will eventually occur. According to Ionson's theory of electrodynamic coupling the efficiency factor f_ϵ at the crossover is about unity. Mullan finds that in the sun f_ϵ ranges from 1/80 to 1/170. Hence f_ϵ at crossover should exceed the solar values by factors of 80 to 170. X-ray data from the Einstein satellite indicate that enhancement of coronal heating efficiencies by factors more than 130 do occur among the early M-dwarfs. He concludes that the greatly enhanced heating efficiency mentioned above is associated with the crossover of t_C and t_A i.e. associated with the resonant behaviour of the electrodynamic coupling efficiency f_ϵ.

Ionson (1985c) extends the work initiated by Heyvaerts and Priest (1984) and comments that the efficiency of DC (nonresonant) mechanisms is severely limited by the constraint that the dissipation time is to be comparable to the correlation time of the convective motions and smaller than the magnetic flux leakage time. He finds that, for solar conditions, DC processes play a dominant role in the heating of young active region loops ($\sim 1.2 \cdot 10^9$ cm) while AC (resonant) mechanisms dominate the heating of active region loops ($\sim 6 \cdot 10^9$ cm) and large scale loops ($\sim 1.6 \cdot 10^{10}$ cm).

Kuperus and Ionson (1985) as an application of LCR circuit approach, demonstrate that the observed large ratio of hard to soft X-ray emission and the bimodal behaviour of black hole accreting X-ray sources such as Cyg X-1 can be described in terms of a magnetically structured accretion disk corona which is electrodynamically coupled to the disk turbulent motions while the disk is thermodynamically coupled to the corona.

Zuccarello et al. (1987) utilize an electric circuit analogy to model the build up and storage of magnetic energy in solar coronal loops. Using the force-free field approximation for a magnetic arcade (whose field lines are sheared by photospheric motions) they demonstrate that the increase of magnetic energy is initially due to the increase of the bulk current I and later mainly due to the increase of the selfinductance L which is determined by the geometry of the magnetic configuration.

Combining basic electric circuit theory with the linearized mhd equations Scheurwater and Kuperus (1988) calculate the input impedance for weakly damped monochromatic Alfvén waves traveling in a magnetoplasma. According to them macroscopic quantities such as stored magnetic and electric energy, dissipated energy, quality factors, damping times etc. can be directly determined from the equivalent circuit representation.

8. Heating by microflares/transients

In this section we deal with steady-state heating of the corona by microflares and transients. Large flares and related phenomena are beyond the scope of the present review.

Priest (1981, 1982) points out that, especially in the strong magnetic field of an active region, the corona is in a state of ceaseless activity and is being heated by many tiny microflarings that are continually generated by the photospheric motion below. The coronal loops that stand out in soft X-ray pictures are those in which most heat is being released and then conducted efficiently along the magnetic field.

Antiochos (1984) presents a model for the solar transition region, $2 \cdot 10^4 < T(K) < 2 \cdot 10^5$, that can account for the persistent and omnipresent redshifts that are observed in the UV emission lines formed at these temperatures. According to his model the redshifted radiation originates from a minority of flux tubes which have higher gas pressures than their surroundings and consequently have their transition regions situated below the chromosphere of their surroundings. The coronal heating in these flux tubes (loops) is impulsive in nature, and this is responsible for the transient mass flows. His study favours theories for coronal heating which involve flare-like magnetic energy release. Flares can be considered as extreme cases of transient coronal heating.

Whitehouse (1985) studies coronal heating and stellar flares. According to him the close correlation of the flare luminosity L_F with the quiescent X-ray luminosity L_X indicates that flares could provide the energy to heat the corona. However, there are obvious differences in the time scales involved. Whereas flares are relatively short-lived phenomena the quiescent X-ray emission is stable over much longer periods. He points out that if the flares themselves are the initial source of energy that heats the corona then the corona must have a large heat capacity that is able to smooth out a fraction of the energy released in stellar flares in the stars considered by him. He argues in favour of coronal heating via magnetic methods.

Antiochos (1987) proposes a coronal heating model which uses the hypothesis that magnetic reconnection acts as a catalyst to initiate the formation of current sheets. His model favours a transient, flare-like coronal heating process. According to him the footpoint motions, by themselves, are unlikely to lead to current sheet formation.

Chiueh and Zweibel (1987) study the general equilibrium structure of current sheets produced by global mhd forces and magnetic reconnection in these sheets. They find that reconnection begins with a linear "sheet-tearing" mode which grows faster than the usual magnetic tearing in a diffuse profile. Their results are encouraging for theories of coronal heating and solar flare models which invoke rapid current dissipation.

Porter and Moore (1987) make an order of magnitude estimate of about 10^4 microflares at any one time and suggest that microflares can supply the necessary energy to heat the corona (Porter et al. 1987, see also Section 10). In fact, they find a global average flux of $5 \cdot 10^5 erg\ cm^{-2}\ s^{-1}$ which is approximately equal to the heating requirement for the quiet corona, lying between the estimates for coronal holes and closed (active) regions (Withbroe and Noyes 1977). The initial energy release takes place in low lying loops spanning small bipoles in the magnetic network. The released energy is transferred to the corona via secondary neutral sheets formed where the expanding low loops push against the legs of adjacent coronal loops. This picture is consistent with the recent work of Machado and Moore (1986) who find that most of the energy in flares is released within closed loops rather than at some interaction site between bipoles but that energy is often transferred out of the flaring loop and into an adjacent structure via interaction sites (Heyvaerts at al. 1977). Thus the corona may be heated by many small variable sources.

Parker (1983b, 1986, 1988, see also Section 6) suggests that the X-ray corona is heated by reconnective dissipation at many small current sheets which are formed all the time as tangential discontinuities between interweaving and winding magnetic filaments. The individual reconnection events are called *nanoflares*. From Si IV and O IV emission line fluctuations observed by Porter et al. (1984) and using the heating flux of $1 \cdot 10^7$ erg cm^{-2} s^{-1} by Withbroe and Noyes (1977), Parker estimates that there are about 5 nanoflares in an area of $4 \cdot 10^{16}$ cm^2 with powers of $8 \cdot 10^{22}$ erg/s and a lifetime of 20 s. The largest nanoflares reach the powers $10^{27} erg/s$ of microflares and produce the isolated turbulent events and high velocity jets observed by Brueckner and Bartoe (1983) (see Section 10). Parker suggests that the observed X-ray corona is simply the superposition of a very large number of nanoflares.

Harrison et al. (1988) investigate the correlations between flaring rates, flare power and quiescent X-ray background of solar active regions and compare them to relations found for dMe stars. The solar relationships are found to be not as clear as those for the dMe stars.

9. Heating by bulk flows or magnetic flux emergence

Hoyle's (1949) hypothesis of coronal heating by gas accretion is no longer believed because of our knowledge of the solar wind. However, it is now generally accepted (see e.g. Vilhu 1987) that mass accretion is a very important heating mechanism for the outer atmosphere of pre-mainsequence stars. The heating theory by spicules or by downflows also involve bulk fluid flow. There is the possibility, however, that such flows are generated by waves and thus might not constitute an independent heating process. Another bulk flow mechanism is the heating by magnetic flux emergence. In this context considerable work has been done on the ejection of plasmoids or coronal bullets (see e.g. Cargill and Pneuman 1986) which, however, will not be covered here as it is not directly relevant for steady chromospheric and coronal heating.

9.1. Heating by spicules

Athay and Holzer (1982) propose that somehow the spicular material is raised well above the height that would be achieved by a projectile of the same initial velocity, thereby obtaining gravitational potential energy much in excess of its initial kinetic energy. This potential energy may then be converted into flow energy as spicular material falls and into internal energy as the flow is slowed, thus providing a source of heat for the solar atmosphere.

From observations they find a global average hydrogen particle flux density carried upward from the network by spicules as

$$F_{sp} = n_{sp}\ v_{sp}\ a_{sp} \approx 1.2 \cdot 10^{15} \quad (cm^{-2}\ s^{-1}) \quad , \tag{9.1}$$

in which a density $n_{sp} \sim 6 \cdot 10^{10}\ cm^{-3}$, a velocity $v_{sp} \sim 20\ km/s$ and a fractional area $a_{sp} \sim 0.01$ have been used. This is about two orders of magnitude larger than the global average solar wind hydrogen flux density of $1.4 \cdot 10^{13} cm^{-2}\ s^{-1}$. Therefore they deduce that most of the material carried upward in spicules must return again to lower layers. The average kinetic temperature in a spicule is $\sim 1.5 \cdot 10^4\ K$.

Again from observations they find an average downward hydrogen flux density of

$$F_{df} = n_{df}\ v_{df}\ a_{df} \approx 1.1 \cdot 10^{15} \quad (cm^{-2}\ s^{-1}) \quad , \tag{9.2}$$

where $n_{df} \sim 6 \cdot 10^9\ cm^{-3}, v_{df} \sim 4\ km/s$ and the fractional network area $a_{df} \sim 0.45$ have been used. The temperature of the downflowing material is found to be in the range $1 - 2 \cdot 10^5\ K$. It is obvious from Eqs. (9.1) and (9.2) that the observed downflow could be the manifestation of the return spicule flow at transition region temperature. Assuming $a_{sp} n_{sp} \tau_{sp} = a_{df} n_{df} \tau_{df}$ they find

$$\frac{\tau_{df}}{\tau_{sp}} = \frac{n_{sp}\ v_{sp}}{n_{df}\ v_{df}} \approx 50 \quad , \tag{9.3}$$

where τ_{sp} is the rise time of spicules and τ_{df} is the downflow time. Thus, if the rise time for the spicule is about a half hour, the downflow time would be about a day (24 hours), which is the average time between spicule ejections at a given location. Therefore they adopt the assumption that the return flow of spicular material occurs almost continuously.

In the light of above observations Athay and Holzer consider a number of possibilities and arrive at the following conclusions. If the spicular material is raised to a height of about $5 \cdot 10^4 km$, spicules can play a significant role in the energy balance of the upper solar atmosphere. Even if no heat is added to the spicular material, compressive heating during its fall can account for the observed downflow of warm material in the network, and the rise and fall of spicules can supply the energy required by radiative losses from the upper chromosphere and transition region in the network. If a substantial amount of energy is supplied as heat to the spicular material (an amount comparable to or greater than that required to lift it), the spicule phenomenon might account for the energy supply required by the corona, transition region and upper chromosphere throughout much of the network. It can, however, provide for the long term maintenance of coronal temperatures only if the magnetic field intensity is about 10 G or less.

The above conclusions do not apply to the maintenance of bright coronal loops, the acceleration of solar wind in coronal holes and other energetically demanding local phenomena. Athay and Holzer emphasize the need for improved understanding of the mechanism whereby spicules are accelerated (and heated) and the need for observations of spicules after they have disappeared in the visual spectrum.

Hollweg (1982b) and Sterling and Hollweg (1988) consider spicules to be produced by acoustic wave propagation in magnetic flux- tubes as discussed in Sec. 2.

Athay (1984) proposes a kinematic model of solar corona and transition region which leads to a different explanation of the origin of spicules. In this model the coronal material collapses and falls creating a mass deficit in the corona which must subsequently be replaced by rising material. This picture is opposite to that proposed by Athay and Holzer (1982). It is assumed that the rise and fall of material occurs in the same flux-tube. This model has two advantages. Firstly it removes the need for a downflow following the spicule upflow. Secondly it offers a natural explanation for a high-temperature downflow and a comparatively low- temperature upflow. This model predicts the existence of "hot spots" in the corona which may result from rapid heating of upflowing material to coronal temperatures.

Schatten and Mayr (1986) suggest that spicules form from the supersonic expansion of material on nearly evacuated flux-tubes embedded within the Sun's convection zone. This allows supersonic but subescape velocities to be attained by the material as it flows outward through the photosphere. In spite of being supersonic, the kinetic energy of spicule material is still insufficient to heat the quiet corona. Through buoyancy changes on evacuated flux tubes, the magnetic field first sucks gaseous matter into the solar atmosphere, then energizes it to form the conventional hot, dynamically expanding, solar corona. This occurs through momentum and energy transport by Alfvén waves and the associated Maxwell stresses concurrently flowing upward with these spicule-type geysers.

9.2. Heating of magnetic flux-tubes by downflows

Hasan and Schüssler (1985) examine, quantitatively, the heating of magnetic flux-tubes in the photosphere and lower chromosphere of the Sun by downflows along the magnetic tube. They perform time-dependent numerical calculations using the magnetohydrodynamic equations, involving (turbulent) viscosity and radiative exchange with the surrounding atmosphere, in the thin tube approximation. They start from a state of hydrostatic and thermal equilibrium. When allowance is made for a downflow initially a transient phase develops which lasts for a few minutes and then a stationary state results that is substantially hotter than the ambient medium over a fairly large height range.

They find that Chapman's facular model can be reproduced remarkably well by adjusting the mass flux entering the tube at the upper boundary. Their results are comparatively insensitive to viscosity ($\leq 10^{12}\ cm^2/s$) while radiative heat exchange is found to be significant. To show the magnitude of heating they give the following net energy gain for the tube: change in enthalpy $8.5 \cdot 10^7\ erg\ cm^{-2}\ s^{-1}$, change in kinetic energy $4 \cdot 10^5\ erg\ cm^{-2}\ s^{-1}$ and change in potential energy $6.3 \cdot 10^7\ erg\ cm^{-2}\ s^{-1}$, which totals to $1.5 \cdot 10^8\ erg\ cm^{-2}\ s^{-1}$. The radiative losses in quiet regions in the lower chromosphere are about $4 \cdot 10^6\ erg\ cm^{-2}\ s^{-1}$ whereas in active regions the corresponding value is $2 \cdot 10^7\ erg\ cm^{-2}\ s^{-1}$ (c.f. Priest 1982).

With a filling factor (fractional area occupied by magnetic structure and downflowing material) of 3% for quiet regions and of 13% for active regions the heating and radiative losses agree with each other quite well. They conclude that the downflows provide an efficient mechanism for heating flux-tubes.

9.3. Heating by magnetic flux emergence

The appearance of solar X-ray bright-points after Gokhale (1975) is associated with the individual emergence of large-scale magnetic flux. For a segment of such an emerging flux rope he gives the following energy estimates: Initial magnetic energy of the segment $\sim 10^{29} erg$, gravitational energy of the residual plasma in the final state $\sim 10^{26} erg$, magnetic energy in the final state $\sim 10^{27} erg$, thermal energy required to heat the plasma to coronal temperatures $\sim 10^{26.5} erg$. Gokhale concludes that the magnetic energy of the segment is sufficient to provide for the radiative and conductive losses $\sim 10^{24.6} erg$ of an X-ray bright-point lasting for ~ 8 hrs.

Uchida and Sakurai (1977) consider heating and reconnection by the interaction of newly emerging magnetic flux tubes with the ambient coronal field. They show that the situation is unstable against the interchange instability. The authors suggest that a continuous relaxation to a lower energy state or a continuous invasion of the emerging magnetic field into the ambient field in the form of fine bundles or thin sheets takes place in a short time scale. The dissipation is assumed to occur in thin current sheets.

Book (1981) considers a bubble-like object, composed of magnetic flux and plasma, emerging from the surface of the Sun. He points out that the exact topology of this object is not critical, e.g. it may have the form of a jet or an arch, and the field lines may remain connected to the solar magnetic dipole. The plasma density ρ is assumed to be less than or equal to that in the surrounding medium, but the initial magnetic pressure $B^2/8\pi$ is assumed to exceed the average ambient pressure. Consequently the structure expands as it rises. This rise continues until pressure balance is reached, or reconnection annihilates the magnetic flux, or the bubble goes out of the top of the atmosphere. Let ΔV be the initial volume of the bubble. Then the work ΔW done by the bubble is given approximately as

$$\Delta W \simeq \frac{B^2}{8\pi} \Delta V \quad . \tag{9.4}$$

Correspondingly a mass $\Delta m = \rho \Delta V$ is injected into the corona. Suppose that structures of this type are continuously generated and the average volume of flux emitted per unit time is $\dot{V}$. Then energy balance implies that

$$\frac{B^2}{8\pi} \dot{V} \simeq (F_{rad} + F_{wind} + F_{cond}) A_\odot \quad . \tag{9.5}$$

where it is assumed that average emitted flux is fully responsible for the heating of the corona. $A_\odot = 4\pi R_\odot^2$ is total surface area of the Sun, $R_\odot$ being the solar radius. F_{rad}, F_{wind}, F_{cond} represent coronal energy losses ($erg\ cm^{-2}\ s^{-1}$) due to radiation, the solar wind flow and thermal conduction, respectively. For the quiet Sun, Eq. (9.5) may be written as

$$\frac{B^2}{8\pi}\ \dot{V} \simeq 3.5 \cdot 10^5\ A_\odot \quad , \tag{9.6}$$

where Withbroe and Noyes' (1977) estimates have been used. Let a fraction f of the Sun's bipolar magnetic regions emit structures with velocity v. Approximately an area of $A_\odot/10$ is occupied by these magnetic regions. Therefore

$$\dot{V} \approx 0.1\ f\ A_\odot\ v \quad . \tag{9.7}$$

If the emerging structure has low density and is propelled by magnetic forces, then v approximately equals the Alfvén speed, v_A. Other mechanisms lead to different estimates for v but the choice is not critical (Book 1981). Assuming $v = v_A$ and an ion number density of $\sim 10^{10}\ cm^{-3}$ one finds using Eqs. (9.6) and (9.7)

$$f \approx 40\ B^{-3} \quad . \tag{9.8}$$

This means that a fraction f, given by Eq. (9.8), of the Sun's magnetically active regions must be emitting flux at any instant of time. The requirement that $f \leq 1$ implies that $B \geq 3\ G$. The upward motion produces a surge which injects mass into the corona at a rate (g/s)

$$\dot{m} = \rho\ \frac{dV}{dt} \approx 1.3 \cdot 10^{-7}\ B^{-2} \quad . \tag{9.9}$$

The requirement that the velocity of the material sprayed out of the solar atmosphere does not exceed the sum of the observed solar wind speed and the escape speed implies $v \leq 10^8\ cm/s$, or $B \leq 50G$. At this upper limit, the total mass injection rate is $\dot{m} \approx 4 \cdot 10^{-11} A_\odot\ g/s$. This is comparable to the observed solar wind mass loss (Withbroe and Noyes 1977).

Adopting a specific model Book (1981) finds that bubbles or other structures containing magnetic flux with field strengths $\sim 20\ G$ and having transverse dimensions like those of the initial stages of bipolar magnetic regions, can supply heat to the solar atmosphere through expansion in amounts sufficient to explain the observed coronal temperatures. If the magnetic forces are responsible for expelling this flux, then bubbles should be emerging from a fraction $f \geq 0.03\%$ of the bipolar magnetic regions at any one time, with characteristic speeds $v_A \leq 500\ km/s$ and lifetimes $R_0/v_A \geq 10^3\ s$ (R_0 being the radius of tube in the adopted model), and should deposit their energy principally within vertical distances $\approx 10^4 - 10^5\ km$ from the surface of the Sun.

Forbes and Priest (1984) have studied reconnection in emerging magnetic flux regions. They find a steady state phase at which reconnection occurs at a slow pace, followed by an impulsive stage. After this there is a quasi-steady state with low level reconnection and subsequently a steady phase with no reconnection.

10. Some relevant observational evidence

Wilson and Bappu (1957) observed that the logarithm of the widths of the Ca II H and K line emission components are linearly correlated with the absolute visual magnitude of the stars of spectral type G0 and later. This relation holds over a very large range of magnitudes.

Leighton et al. (1962) reported the observation of an oscillation in the photosphere with a period $296 \pm 3\ s$, the five-minute oscillation, with a velocity amplitude of about $0.4\ km/s$.

Wilson (1963) finds that the average intensity of the Ca II H and K emission for main-sequence stars of types G0-K2 in the Hyades, Praesepe, Coma, and Pleiades clusters is higher than for similar local field stars. He suggests that the intensity of the H-K emission is related to the stellar magnetic field and that the decrease of the chromospheric emission is due to the decrease of the magnetic field with age.

Belcher and Davis (1971) presented observational results of Alfvén (body) waves in the interplanetary medium, having wavelengths in the range 10^3 to $5 \cdot 10^6\ km$.

Skumanich (1972) found that the decrease of the Ca II H and K emission with age is closely correlated with the decrease of the stellar rotation and the Li abundance.

Boland et al. (1973), from their observations of spectral line profiles emitted by ions in the temperature range $10^4 - 10^5\ K$, show the existence of a non-thermal kinetic energy component in the chromosphere - corona transition region. They find a mean mechanical energy flux of $5 \cdot 10^5\ erg\ cm^{-2}\ s^{-1}$ and conclude that the corona is heated by mechanical energy carried from lower layers by propagating sound waves.

Beckers and Artzner (1974) describe the properties of dark structures, seen in the K-line wings, which seem to propagate upward in the solar atmosphere. They point out that these structures may be related to shocks that heat the chromosphere and corona.

Deubner (1975), on the basis of his observations, concludes that the five-minute ($300\ s$) oscillations may in fact be interpreted as low wavenumber non-radial acoustic eigenmodes of the subphotospheric layers of the solar atmosphere.

Observations of Fossat and Ricort (1975) show an oscillation of $180\ s$ period, the three minute oscillation, in the chromosphere. Their velocity oscillation amplitudes indicate that shock waves begin to appear in the low chromosphere. They point out that chromospheric acoustic flux is found to be in good agreement with the value needed for heating the chromosphere.

Giovanelli (1975) reports observations of Alfvén waves propagating along H_α fibrils over disturbed and quiet supergranules and in a young system of superpenumbral fibrils associated with a small, young sunspot. According to him, the Alfvén waves propagate with a velocity of $60 - 70\ km/s$ along the fibrils in outward direction.

Vernazza et al. (1975) analyze the time structure of the intensity of solar chromospheric and coronal EUV lines. They find evidence for the presence of shock waves in the chromosphere and the transition region and suggest that the solar chromosphere and corona are heated by non-periodic waves.

Beckers (1976), assuming the horizontal velocities observed in the sunspot photosphere to be caused by Alfvén waves, derives an Alfvén flux of $\sim 10^{10}$ $erg\ cm^{-2}\ s^{-1}$. Beckers and Schneeberger (1977), from the line-width in coronal arches above sunspots, estimate the amount of Alfvén wave flux escaping from the sunspot into the solar corona and find it to be less than $4 \cdot 10^7$ $erg\ cm^{-2}\ s^{-1}$.

Deubner (1976) obtains observational evidence of acoustic waves in the solar atmosphere through an analysis of power spectra of velocity fluctuations derived from high spatial resolution solar spectra. He finds that there is considerable energy ($\sim 10^8 - 10^9$ $erg\ cm^{-2}\ s^{-1}$) in the frequency range below 0.3 Hz (> 20 s period) in spectral lines representing the solar atmosphere from the visible surface to the height of formation of H_α. A critical analysis by Cram et al. (1979), Durrant (1980) and Deubner et al. (1982) shows that this acoustic flux has been overestimated.

Foukal (1976), from a study of EUV emission of the 22 largest observed sunspots, finds that a heat input of about 10^{-4} $erg\ cm^{-3}\ s^{-1}$ is required along the full length of a loop and it is largest off the loop axis. An analysis of energy and pressure balance suggests that plasma is falling under gravity down both sides of the loop.

Doschek and Feldman (1977) report observations of some coronal lines over a quiet sun region, a coronal hole and two active regions. They find that iron lines ($\sim 10^6$ K) are not observed in the coronal hole indicating that most of the plasma there is at a lower temperature than 10^6 K. The nonthermal velocities in the coronal hole and quiet regions are ~ 20 km/s but in active regions they are substantially less.

Levine and Withbroe (1977) present observations of sudden mass evacuations in active region loops. They point out that the cool cores of loops are likely to be no more than a few hundred kilometers in radius and that several such cool threads may be imbedded in a common hot outer sheath. Their observations show evidence for continual energy input to the loop.

Nolte et al. (1977) examine observed reconfigurations of coronal X-ray and XUV emitting structures. They point out that reconnection occurs in regions much smaller than telescopic resolution and existing observations are generally not sufficient to show in detail how much reconnection has occurred.

Withbroe and Noyes (1977) review the observations of mass and energy flow in the solar chromosphere and corona. Table II gives a summary of chromospheric and coronal heating fluxes.

Athay and White (1978, 1979) observed solar oscillations, via the OSO 8 satellite, in the middle chromosphere and find an estimated energy flux in sound waves of approximately $1 \cdot 10^4$ $erg\ cm^{-2}\ s^{-1}$ which is much below the heat input needed for the upper chromosphere and corona.

Bonnet (1978) reviews observational results and reports the existence of downflow velocities as large as 22 km/s in the transition layer and of short period waves of period 95 s in the chromosphere in addition to the well-known 300 s and 180 s photospheric and chromospheric oscillations.

Gerassimenko et al. (1978) observed X-ray emission from three active region loops. They find evidence for continuous (or intermittent) heating.

Giovanelli et al. (1978) observed motions in solar magnetic tubes. According to them the disturbance velocities never exceed about 0.1 of the velocity of sound,

TABLE II
Chromospheric and coronal fluxes after Withbroe and Noyes (1977)

	Quiet sun	Coronal hole	Active region
Transition layer pressure (dyn cm^{-2}):	$2 \cdot 10^{-1}$	$7 \cdot 10^{-2}$	2
Coronal fluxes (erg cm^{-2} s^{-1}):			
conductive	$2 \cdot 10^5$	$6 \cdot 10^4$	10^5 to 10^7
radiative	10^5	10^4	$5 \cdot 10^6$
wind	$\leq 5 \cdot 10^4$	$7 \cdot 10^5$	$< 10^5$
total	$3 \cdot 10^5$	$8 \cdot 10^5$	10^7
Chromospheric flux (erg cm^{-2} s^{-1}):			
total	$4 \cdot 10^6$	$4 \cdot 10^6$	$2 \cdot 10^7$

suggesting the absence of large-amplitude shock waves anywhere in the magnetic or non- magnetic solar atmosphere between the levels of formation of Fe I 5166Å and H_α.

Linsky and Ayres (1978) develop a method for estimating the non- radiative heating of stellar chromospheres by measuring the net radiative losses in strong Fraunhofer line cores, and apply the method to observation of the Mg II resonance lines in a sample of 32 stars including the sun. They find a small dependence of chromospheric nonradiative heating on stellar surface gravity, contrary to the large effect predicted by calculations of the acoustic flux generated in the convection zones.

Basri and Linsky (1979) present IUE high resolution Mg II line spectra of 15 stars of spectral type G2 - M2. They do not find a gravity dependence of the Mg II emission.

Cheng et al. (1979) study line profiles of some forbidden lines in quiet and active coronal regions. From line broadening they find a non-thermal mass-motion velocity of $10 - 25\ km/s$ at $1.7 \cdot 10^6\ K$ (Fe XII), of $10 - 17\ km/s$ at $1.5 \cdot 10^6\ K$ (Fe XI) and $10 - 20\ km/s$ at $9.3 \cdot 10^5\ K$ (Si VIII). The intensities of lines in active regions are an order of magnitude greater than those in quiet regions. They discuss their results in the light of proposed heating mechanisms and find that neither acoustic waves nor mhd waves are able to explain their observations satisfactorily.

Nicolas et al. (1979) investigate the energy balance and the pressure in the transition region for a number of network and active region features observationally and analytically. They find that the energy balance requires the existence of some source of energy for which they assume dissipation of turbulent motions. For this heating rate ($erg\ cm^{-3}\ s^{-1}$) they take

$$\epsilon_t = \rho \frac{(\Delta v)^3}{\Delta x} \quad , \tag{10.1}$$

where Δv represents turbulent or non-thermal velocity and Δx is the scale length in which the turbulence is dissipated.

Provost and Mein (1979) interpret the small phase lag between velocities observed at different chromospheric levels as being due to acoustic waves reflected by the very hot atmospheric layers of the chromosphere - corona transition zone. They infer that in upper chromospheric layers magnetoacoustic waves become important.

Schmieder (1979) made observations of the Mg I line at 5172.7 Å . She finds that radiative dissipation seems to be efficient up to an altitude of 600 km, in the evanescent wave range ($3-5\ mHz$). Some energy can be transported but in the acoustic wave range ($5-8\ mHz$) no energy propagation to the high chromosphere and corona is detected. It has been pointed out that the presence of pure downward propagating waves around the temperature minimum can be anticipated.

Brown and Harrison (1980) report observations of disk-center solar continuum brightness fluctuations and they interpret it as evidence for internal gravity waves trapped in the solar photosphere and chromosphere.

Brueckner (1980) presents a high resolution view of the solar chromosphere and corona. He reports some explosive events which have a rise time of less than 20 seconds. According to him supersonic motions are confined to small elements in the transition zone and the supersonic motions in the corona are observed as intensity spikes in the ultraviolet spectrum. He points out that spikes may be parts of high loops in a coronal hole and concludes that the heating of the solar atmosphere as well as the energy needed for the propulsion of the solar wind, is caused by magnetoacoustic waves or electric currents.

From X-ray observations and magnetic field data Golub et al. (1980, see also Hollweg and Sterling 1984) obtain quantitative information of the physical properties of coronal loops. The observations indicate a relation $p = 3.0 \cdot 10^{-3}\ B^{1.6}$ between the gas pressure and the magnetic field strength of the loops. For their active region loops they find typical gas pressures of $p = 1.5\ dyn/cm^2$, magnetic field strengths $B = 50\ G$ and loop lengths of $l_{||} = 2 \cdot 10^5\ km$. The loops with a diameter $l_\perp$ show an aspect ratio $l_\perp/l_{||} = 0.1 - 0.4$.

Kohl et al. (1980) measured coronal kinetic temperatures in a coronal hole and quiet region. In the quiet region, they find the temperature to fall from about $2.5 \cdot 10^6\ K$ at $2\ R_\odot$ to about $10^6\ K$ at $\sim 4\ R_\odot$ whereas in the coronal hole at $2.5\ R_\odot$ the temperature is $1.8 \cdot 10^6\ K$.

Schmieder and Mein (1980) determine the mechanical flux integrated over the frequency range $\nu < 10\ mHz$ to be $2 \cdot 10^3\ erg\ cm^{-2}\ s^{-1}$ in the middle chromosphere. They point out that this flux is not sufficient to balance the energy losses of the transition layer and corona.

Stencel et al. (1980) present high resolution spectra of the emission cores of the Mg II resonance doublet at 2800 Å in a selection of 54 stars covering a range of spectral type from F8 to M5. Their results indicate that at a given T_{eff} ($< 5000\ K$) the value of the normalized radiative flux, $F(MgII)/\sigma T_{eff}^4$ for supergiants is about 3-4 times larger than that for giants, corresponding to an approximate gravity dependence $F(MgII)/\sigma T_{eff}^4 \sim g^{-0.25}$. In spite of appreciable scatter in the data, the trend for the temperature dependence, suggested by these results, is approximately $F(MgII)/\sigma T_{eff}^4 \sim T_{eff}^4$ (see, e.g. Cassinelli and MacGregor 1986).

Brueckner (1981) describes the dynamics of active regions on the basis of his 1975-1978 observations. According to him observations do not show permanent upflows except during explosive events. He points out that if the observed transition region downflows over active regions (20-60 km/s) and sunspots (up to 150 km/s) are free- falling material then the existence of loops with an altitude range from 750 km to 40000 km must be assumed. He remarks that explosive events may be considered as the source of material being propulsed upward which may come down as downflows.

Bruner (1981), through the analysis of OSO 8 observations, concludes that the upper limit to the acoustic flux available for coronal heating, is three orders of magnitude too low to heat the corona.

Mein and Schmieder (1981) report that the mechanical flux for waves having periods in the range 400 and 120 s decreases from $8 \cdot 10^7$ $erg\ cm^{-2}\ s^{-1}$ in the photosphere (at height $\sim$ 170 km) to $2 \cdot 10^3$ $erg\ cm^{-2}\ s^{-1}$ (at $\sim$ 1500 km) and point out that heating of the corona by shock wave dissipation seems improbable. Mein (1981) includes acoustic waves from 60 to 30 s and finds the total flux in all frequencies to be about $4 \cdot 10^3$ $erg\ cm^{-2}\ s^{-1}$ which is about two orders of magnitude less than required for the heating of the transition region and corona.

Cram and Damé (1983), through their high spatial and temporal resolution observations of the solar Ca II H line, show that the three-minute chromospheric oscillation involves upward-propagating excitation which leads to intense heating in the cell points. They suggest that the three-minute oscillation is responsible for heating the quiet chromosphere lying outside network regions.

Giovanelli and Beckers (1983) report simultaneous observations in the H_α and K lines of motions along H_α fibrils. They find that close to the network, the velocity of propagation is of the order 12 km/s towards or away from the network; further away the patterns propagate away from the network with velocities of the order of 75 km/s. The latter are interpreted as Alfvén waves, the former as due most likely to variations in the longitudinal velocities along the fibrils.

Koutchmy et al. (1983) infer the existence of Alfvén waves of period 84.5 and 43 s from their observations. They find a total flux $\sim 4 \cdot 10^5$ $erg\ cm^{-2}\ s^{-1}$ in the two harmonics.

Brueckner and Bartoe (1983) observe high-energy events ('turbulent events' and 'high-velocity jets') in the corona above the quiet sun. Turbulent events show high turbulence (up to 250 km/s), are confined to small areas ($<$ 1500 km) and have average lifetimes of 40 s. The average energy of a turbulent event is $7 \cdot 10^{23}$ erg. With 753 events per sec one has a total power of $6 \cdot 10^{26}$ erg/s and a heating flux of $9 \cdot 10^3$ $erg\ cm^{-2}\ s^{-1}$ for the whole sun in turbulent events. High-velocity jets have moving material with velocities ($\approx$ 400 km/s) exceeding the sound speed ($\approx$ 120 km/s) and are confined to areas $<$ 3000 km. A single jet with an energy of $3 \cdot 10^{26}$ erg carries a mass of $3 \cdot 10^{11}$ g to an altitude of $4000 - 16000$ km and has a maximum lifetime of 80 s. With 24 jets per second one finds a power of $7 \cdot 10^{27}$ erg/s and a heating flux of $1 \cdot 10^5$ $erg\ cm^{-2}\ s^{-1}$ over the whole sun.

Cook et al. (1984) report observations of small loop- or prominence- like events which exhibit rotational velocities of approximately 50 km/s.

Lindsey and Kaminski (1984) present observations of local intensity variations in the 300 - 800 μm solar continuum. They find significant enhancement in spectral power for frequencies between 3 and 7 mHz and attribute these and higher frequency variations to the adiabatic response of the chromospheric medium to compression waves.

The photoelectric observations of chromospheric sunspot oscillations of Lites (1984) suggest that the oscillations are upward propagating acoustic (or slow mode) disturbances and that they become nonlinear and develop into shock waves in the upper layers. They cause a significant increase in the radiative output of the umbral chromosphere, indicating the possibility of nonthermal heating at these levels. Oscillations are present in the outer regions of penumbrae with frequencies and phase relationships that suggest the possibility of magnetogravity waves.

Pasachoff and Landman (1984) report observations and an analysis of the 1980 solar eclipse without correction for atmospheric and instrumental contributions. They detected excess power between 0.5 and 2 Hz and suggest that these oscillations could be associated with Alfvén waves (surface waves) that are trapped on loops a few thousand kilometers long or with fast-mode waves that are trapped on loops a few thousand kilometers in diameter.

Porter et al. (1984) report observations in the ultraviolet of sites of enhanced intensity within an active region on the sun. Their results imply heating due to magnetic field reconnection taking place almost stochastically. In addition to this they find that the events involving only a modest energy release occur most frequently.

Schrijver et al. (1984) present soft X-ray spectra and an analysis for 34 late type stars. They find two types of coronal loop structures: one with a high temperature ($2 \cdot 10^7\ K$), most pronounced on giants, and the other with a low temperature ($2 \cdot 10^6\ K$), dominating the emission of dwarfs.

Fourier analyzing high spatial resolution Doppler shift observations in solar spectral lines Deubner (1985) finds that the phase spectra are seriously contaminated by coherent spurious signals due to seeing and instrumental misguiding at frequencies above $\sim 10\ mHz$. Correcting the spectra he detects genuine signals of acoustic waves up to frequencies of $\sim 25\ mHz$ ($P = 40\ s$).

Habbal et al. (1985) analyze Skylab EUV observations of an active region near the solar limb. They remark that one coronal heating mechanism that can account for the observed behaviour of the EUV emission from McMath region 12634 is the dissipation of fast mode mhd waves.

Martens et al. (1985) report X-ray observations by the Hard X-ray Imaging Spectrometer (HXIS) of the thermal evolution of a loop for about 15 hours. According to them the loop was heated steadily at a rather large rate of 0.6 $erg\ cm^{-3}\ s^{-1}$. They interpret this heating as due to the dissipation of magnetic fields in thin current sheets via ion-kinetic tearing (Galeev et al. 1981) which leads to enhanced resistivity.

Withbroe et al. (1985) report results of observation of coronal temperatures, heating and energy flow in the polar region of the sun at solar maximum. They find the temperatures to decrease from $1.2 \cdot 10^6\ K$ at 1.5 $R_\odot$ to $6 \cdot 10^5\ K$ at 4 $R_\odot$. Their estimated fluxes at 1.5 $R_\odot$ are exhibited in Table III.

Foukal et al. (1986) analyze limb spectra of five post flare loops and three active prominences. They find a DC electric field, of order 100 V/cm , whose projection in the sky plane lies transverse to the loop magnetic field vector. They suggest that such a field might arise from reconnection in post- flare loops.

TABLE III
Coronal energy fluxes ($erg\ cm^{-2}\ s^{-1}$) at $1.5R_\odot$

Kinetic	4.9
Enthalpy	$5.8 \cdot 10^3$
Gravitational	$-1.5 \cdot 10^4$
Conductive	$1.0 \cdot 10^4$
Mechanical	$1.1 \cdot 10^3$
MHD wave	$7.2 \cdot 10^3$

Advances in the detector technology permits spectroscopy of the Sun in the $10 - 12\mu m$ infrared window (Deming 1987). The large Zeeman sensitivity of the emission lines makes the infrared of special interest for the study of magnetic fields. Time series observations of OH lines showed acoustic energy near $3 - 4.5 mHz$.

Harrison (1987) reports solar soft X-ray (3.5-5.5 keV) pulsations of 24 min period using HXIS data from a very compact active region. According to him, the pulsation is caused either by fast- or Alfvén- modes or by Alfvénic surface waves.

Lindsey and Roellig (1987) used the NASA Infrared Telescope Facility to observe local continuum brightness variations of the quiet sun at 350 and 800 μm simultaneously. They find a phase lag in the variations of the two bands and attribute it to the non-adiabatic response of the chromosphere to compression which could be an important mechanism for the dissipation of mechanical energy.

Pasachoff and Ladd (1987) report observations of the 1983 eclipse for high frequency oscillations in the corona. Their results indicate the existence of power in coronal loops at high frequencies, suggesting heating by fast mode waves.

Porter et al. (1987) observed sequences of images of the quiet sun in the C IV 1548 Å line with the Ultraviolet Spectrometer and Polarimeter on the SMM satellite. These images reveal localized brightenings throughout the magnetic network. Some bright sites are short lived, lasting less than 2 minutes, which is the repeat time of the observations, while others persist through several rasters and sometimes for an hour. By comparison with Kitt Peak magnetograms they find that the sites lie on neutral lines of small magnetic bipoles, with most of the long-lived sites corresponding to the stronger bipoles. After comparison with He I 10830 Å spectroheliograms it emerges that the long-lived C IV sites lie on He I dark points, corresponding to X-ray bright points. From plots of intensity fluctuations they find that the enhancements at both short- and long-lived sites are the result of localized impulsive heating events, occurring intermittently at the short-lived sites and in more rapid succession at the long-lived sites. The number of these events and their visibility in the wings of C IV are found to be consistent with their identification as the explosive events seen in UV spectra obtained by the NRL High-Resolution Telescope and Spectrograph. The location of these events on neutral lines, together with their rapid rise times, short durations, and small spatial scales (< 6") suggests that they are microflares occurring in short (< 4000 km) magnetic loops. Porter et al. conclude that their observations favor microflares to be the source of both

TABLE IV

Temperature T, electron density n_e, gas pressure p, density ρ, magnetic field strength B, Alfvén speed v_A, nonradial area factor f, and mechanical heating rate ϵ_H for quiet regions and equatorial coronal holes at a distance $R = 1.1 R_\odot$ after Withbroe (1988).

	$T\ (K)$	$n_e\ (\frac{1}{cm^3})$	$p\ (\frac{dyn}{cm^2})$	$\rho\ (\frac{g}{cm^3})$	$B\ (G)$	$v_A\ (\frac{km}{s})$	f	$\epsilon_H (\frac{erg}{cm^3 s})$
$q.r.$	$1.5 \cdot 10^6$	$1.3 \cdot 10^8$	$5.2 \cdot 10^{-2}$	$2.5 \cdot 10^{-16}$	8	$1.4 \cdot 10^3$	1.5	$4.6 \cdot 10^{-5}$
$e.h.$	$1.2 \cdot 10^6$	$4.7 \cdot 10^7$	$1.5 \cdot 10^{-2}$	$9.0 \cdot 10^{-17}$	9	$2.6 \cdot 10^3$	1.3	$9.5 \cdot 10^{-6}$

spicules and the heating of the corona.

Schrijver (1987b) analyzes magnetograms (KTPN) and UV spectroheliograms (ATM spectrometer) of solar active regions. He finds that the magnetic field, averaged over the magnetic plage of active regions, has a nearly constant value $100 \pm 20\ G$ whereas the mean magnetic field in quiet regions is of the order of $10\ G$. He points out that the observed radiative losses from the corona over active regions agree with quantitative predictions of heating by shearing of the magnetic field by granulation. From Schrijver's X-ray flux - magnetic flux correlation a relation $F_x(erg\ cm^{-2}\ s^{-1}) = 4.0 \cdot 10^4\ B(G)$ has been derived (see Vilhu 1987) which allows to estimate the X-ray fluxes of stars with large filling factor f. Schrijver et al. (1989) find a similar relation $\Delta F_{CaII} = 0.6 \log < fB > +4.8$ between the CaII excess flux and the magnetic flux.

Shore et al. (1987) present observations of rapidly oscillating Ap stars and conclude that their observations rule out Alfvén wave heating in these stars.

Deubner et al. (1988) report results of observation on the energy of short period acoustic waves in the solar atmosphere. At a height of $\sim 300\ km$ in the atmosphere they find a mechanical energy flux of $\sim 1.0 \cdot 10^7\ erg\ cm^{-2}\ s^{-1}$. They point out that the amount of energy carried by short period acoustic waves is particularly uncertain in the range of frequencies ($\geq 10\ mHz$) which is potentially relevant for acoustic heating.

Withbroe (1988) presents models of radial variations of coronal temperatures, densities and outflow speeds up to $10\ R_\odot$ in several types of coronal holes and a quiet region of the corona using data from remote sensing and in situ instruments as empirical constraints. He finds a total nonradiative heating flux in quiet regions of $F_M = 8 \cdot 10^5\ erg\ cm^{-2}\ s^{-1}$ with a damping length of $L_M = 0.25 R_\odot$ and in an equatorial coronal hole of $F_M = 4.2 \cdot 10^5\ erg\ cm^{-2}\ s^{-1}$ with $L_M = 0.63 R_\odot$. He points out that most of the energy which heats the coronal plasma appears to be dissipated within $2\ R_\odot$ of the solar surface. In the following Table IV we list several typical values found by Withbroe for quiet regions and equatorial coronal holes at $R = 1.1 R_\odot$. For the magnetic field a mean field of $B = 6 \cdot 10^{-5}\ G$ at 1 AU was taken, and $Bf(R)R^2 = const.$, where f is the nonradial areal expansion factor given by Withbroe. At 1 AU, $f = 5$ was assumed. The mechanical heating rates were computed using $\epsilon_H = F_M / L_M$. Withbroe reports nonthermal velocities of $\delta v \approx 20\ km/s$ at $R = 1.1 R_\odot$ obtained from FeX line broadening measurements which are consistent with radio scintillation measurements obtained at $10 R_\odot$.

Anderson and Athay (1989a,b) determined the chromospheric heating flux requirements by computing semiempirical models which include mechanical heating. Taking into account Fe II, Ca II, Mg II and H cooling they reproduce the temperature structure determined by Vernazza et al. (1981) derived from continuum observations and find that a mechanical heating flux of $F_M = 1.4 \cdot 10^7\ erg\ cm^{-2}\ s^{-1}$ is needed. The decrease of this flux with height is consistent with the heating by acoustic waves.

Fleck and Deubner (1989) as well as Deubner and Fleck (1989) studied time series of brightness and velocity fluctuations measured simultaneously in three chromospheric lines (Ca II 8542, 8498, Fe I 8496) both in the chromospheric network and the supergranulation cell interior. They find that the 3-min oscillation is a standing wave phenomenon with a 90^o phase shift between brightness and velocity oscillations and very large phase speed. The 3-min oscillation can be seen only in the dark supergranulation cell interior and not in the network where the 5-min oscillation dominates. The short period acoustic wave flux is seen to propagate up to a *magic height* of $800 - 1000\ km$, above this height the short period waves seem to be standing, exhibiting very large phase speed.

Von Uexküll et al. (1989) investigate the dynamics of the chromospheric network at supergranulation cell boundaries employing two-dimensional H_α-line spectroscopy. While in the cell interior there are oscillatory motions consistent with acoustic waves, the cell boundaries show stochastic motions with rms velocities of $4\ km/s$. With lifetimes of the H_α fine mottles of $t = 400\ s$ the authors estimate the heating flux entering the network due to photospheric granular motions using Eq. (6.3). For loops with $B = 50\ G$ and $l_{||} = 4 \cdot 10^3,\ 1 \cdot 10^5\ km$ they find heating fluxes of $F_H = 3.2 \cdot 10^7,\ 1.6 \cdot 10^6\ erg\ cm^{-2}\ s^{-1}$, respectively, which are sufficient for the energy balance of the high chromosphere.

11. Conclusions

Stellar chromospheres and coronae are heated by a multitude of viable mechanisms. However, due to both the insufficient observational evidence and the relatively poor theoretical development of most heating processes, it is not possible at the present time to identify with certainty the dominant heating mechanism in a given situation.

1. In late-type stars short period acoustic waves are thought to contribute a basic but weak background heating which is independent of rotation. These waves are generated by the turbulent motions of the stellar surface convection zones. Short period ($\sim P_A/10$, where for P_A see Eq. 2.10) waves heat the chromosphere by shock dissipation and in regions of significant hydrogen ionization by ionization pumping. The chromospheres of very late red giants with their low rotation rates seem to be dominated by acoustic heating.

2. The solar 3- and 5- min oscillations with their large phase shifts are standing waves and do not dissipate via shocks in the chromosphere. In late-type giant stars this acoustic wave type, however, contributes to heating as e.g. in Mira stars and to mass loss.

3. Coronal structures in early-type stars are heated by strong acoustic shocks which are generated out of small disturbances by radiative amplification. Conse-

quently the heating in early-type stars depends on the bolometric luminosity and not on rotation.

4. The dominant contribution to the chromospheric and especially coronal heating in rotating late-type stars is the magnetic field-related heating which shows strong dependence on rotation. The magnetic heating mechanisms are direct current dissipation of fields and MHD-wave dissipation. For the latter, magnetoacoustic waves (fast- and slow-mode waves in homogeneous fields and acoustic-like longitudinal waves in flux-tubes) as well as shear (transverse and torsional) Alfvén waves have been considered. In structured magnetic fields in addition to body waves, surface waves are present. The waves are generated by the turbulent and convective motions in the stellar surface convection zones. Their energy is transmitted via foot point motions into the photospheric, chromospheric and coronal magnetic structures. Coronal loops gain energy by convective foot point motions and differential rotation. In addition energy enters the outer stellar atmospheres by magnetic flux emergence.

5. Fast-mode waves can heat coronal loops provided they are produced locally in the corona. Those produced below the transition layer due to strong refraction are not able to reach the corona. In the corona, fast-mode waves may be generated by mode- coupling.

6. Slow-mode waves and longitudinal tube waves propagate without much difficulty from the convection zone to the chromosphere but suffer radiation damping. Due to their acoustic wave nature they readily heat the low and middle chromosphere through shock- dissipation. These waves moreover are easily generated by non-linear mode-coupling from other wave types.

7. Transverse and torsional Alfvén body waves easily travel to coronal heights and heat coronal loops by resonant heating, phase-mixing, mode-coupling, turbulent heating and at great heights by Landau damping.

8. Alfvénic surface waves heat coronal loops by resonant absorption in thin sheaths.

9. Coronal holes are heated by the dissipation of shear Alfvén waves by phase-mixing, mode-coupling and turbulent heating. The heating provides the thermal energy that drives the solar wind. Non-dissipated Alfvén waves contribute their momentum to the solar wind, especially to high speed streams, boosting their velocity above that available from thermal expansion alone.

10. Coronal loops may also be heated by anomalous current dissipation or magnetic reconnection. These mechanisms seem to provide intermittent heating unless some external mechanism maintains favourable conditions constantly. The heating occurs in very thin sheaths or layers. Complex magnetic reconnection and turbulent heating at vortex structures seem quite promising.

11. The resonant LCR circuit approach provides a simple way to obtain results which could be verified by observations more readily. It also provides a formalism to unify various proposed mechanisms.

12. Microflares/transients heat the corona. The basic mechanism of their production appears to be the same as by direct current dissipation.

13. Large scale flows like spicular flows may ultimately be due to heating by other means. Magnetic flux emergence as seen in coronal bright points introduces

considerable energy into the chromosphere and corona. The important heating by mass accretion observed in pre-main sequence and protostars is not covered in this work.

14. Observations of photospheric, chromospheric and coronal velocities, temperatures and magnetic fields etc. on the scales required by the various proposed mechanisms are seriously lacking. Instruments which are capable to resolve the important spatial scales of the heating mechanisms are absolutely essential. Unless these vital observations are available, discrimination among the different heating mechanisms and specification of their atmospheric domain of validity remains speculative. On the other hand a more detailed theoretical modelling of most heating mechanisms is also urgently needed.

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References

Abdelatif, T.E.: 1987, Astrophys. J. **322**, 494
Alfvén, H.: 1947, Mon. Not. Roy. Astr. Soc. **107**, 211
Anderson, L.S., Athay, R.G..: 1989a, Astrophys. J. **336**, 1089
Anderson, L.S., Athay, R.G..: 1989b, Astrophys. J., in press
Antiochos, S.K.: 1984, Astrophys. J. **280**, 416
Antiochos, S.K.: 1987, Astrophys. J. **312**, 886
Anzer, U., Galloway, D.J.: 1983, in Solar and Stellar Magnetic Fields, IAU Symp. **102**, J.O. Stenflo Ed., p. 339
Assis, A.S., Busnardo-Neto, J.: 1987, Astrophys. J. **323**, 399
Athay, R.G.: 1976, The Solar Chromosphere and Corona: Quiet Sun, Reidel, Dordrecht, Holland
Athay, R.G.: 1984, Astrophys. J. **287**, 412
Athay, R.G., White, O.R.: 1978, Astrophys. J. **226**, 1135
Athay, R.G., White, O.R.: 1979, Astrophys. J. **229**, 1147
Athay, R.G., Holzer, T.E.: 1982, Astrophys. J. **255**, 743
Avrett, E.H.: 1981, in Solar Phenomena in Stars and Stellar Systems, R.M. Bonnet, A.K. Dupree Eds., Reidel, Dordrecht, Holland, p. 173
Barnes, A.: 1966, Phys. Fluids **9**, 1483
Barnes, A.: 1968, Astrophys. J. **154**, 751
Basri, G.S., Linsky, J.L.: 1979, Astrophys. J. **234**, 1023
Beckers, J.M.: 1976, Astrophys. J. **203**, 739
Beckers, J.M., Artzner, G.: 1974, Solar Phys. **37**, 309
Beckers, J.M., Schneeberger, T.J.: 1977, Astrophys. J. **215**, 356

Belcher, J.W., Davis, L. Jr.: 1971, J. Geophys. Res. **76**, 3534
Biermann, L.: 1946, Naturwiss. **33**, 118
Biermann, L.: 1948, Z. Astrophys. **25**, 161
Billings, D.E.: 1966, A Guide to the Solar Corona, Academic Press, New York, USA
Bird, G.A.: 1964a, Astrophys. J. **139**, 675
Bird, G.A.: 1964b, Astrophys. J. **139**, 684
Bird, G.A.: 1965, Astrophys. J. **141**, 1455
Blake, M.L., Sturrock, P.A.: 1985, Astrophys. J. **290**, 359
Bobrova, N.A., Syrovatskii, S.I.: 1979, Solar Phys. **61**, 379
Bodo, G., Ferrari, A., Massaglia, S., Rosner, R., Vaiana, G.S.: 1985, Astrophys. J. **291**, 758
Bodo, G., Ferrari, A., Massaglia, S., Rosner, R.: 1987, Astrophys. J. **313**, 432
Bohn, H.U.: 1981, Ph.D. Thesis, Univ. Würzburg, Germany
Bohn, H.U.: 1984, Astron. Astrophys. **136**, 338
Boland, B.C., Engstrom, S.F.T., Jones, B.B., Wilson, R.: 1973, Astron. Astrophys. **22**, 161
Bonnet, R.M.: 1978, Space Sci. Rev. **21**, 379
Book, D.L.: 1981, Comments Plasma Phys. Controlled Fusion **6**, 193
Braginskii, S.I.: 1965, Rev. Plasma Phys. **1**, 205
Bray, R.J., Loughhead, R.E.: 1974, The Solar Chromosphere, Chapman and Hall, London, England
Brinkley, S.R. Jr., Kirkwood, J.G.: 1947, Phys. Rev. **71**, 606
Brown, T.M., Harrison, R.L.: 1980, Astrophys. J. **236**, L169
Browning, P.K.: 1984, in the Hydromagnetics of the Sun, Proceedings of the Fourth European Meeting on Solar Physics, T.D. Guyenne, J.J. Hunt Eds., ESA SP-220, p. 129
Browning, P.K., Priest, E.R.: 1984, Astron. Astrophys. **131**, 283
Browning, P.K., Priest, E.R.: 1986, Astron. Astrophys. **159**, 129
Browning, P.K., Sakurai, T., Priest, E.R.: 1986, Astron. Astrophys. **158**, 217
Brueckner, G.E.: 1980, in Highlights of Astronomy **5**, P.A. Wayman Ed., Reidel, Dordrecht, Holland, p. 557
Brueckner, G.E.: 1981, in Solar Active Regions, F.Q. Orrall Ed., Colorado Ass. Univ. Press, Boulder CO, USA, p. 113
Brueckner, G.E., Bartoe, J.D.F.: 1983, Astrophys. J. **272**, 329
Bruner, E.C. Jr.: 1981, Astrophys. J. **247**, 317
Cargill, P.J., Priest, E.R.: 1983, Astrophys. J. **266**, 383
Cargill, P.J., Pneuman, G.W.: 1986, Astrophys. J. **307**, 820
Cassinelli, J.P.: 1979, Ann. Rev. Astron. Astrophys. **17**, 275
Cassinelli, J.P., MacGregor, K.B.: 1986, Physics of the Sun, Vol. III, P.A. Sturrock, T.E. Holzer, D. Mihalas, R. Ulrich Eds., Reidel, Dordrecht, Holland, p. 70
Castor, J.I.: 1987, in Instabilities in Luminous Early Type Stars, H.J.G.L.M. Lamers, C.W.H. de Loore Eds., Reidel, Dordrecht, Holland, p.159
Chen, F.F.: 1974, Introduction to Plasma Physics. Plenum Press, New York, USA, p. 39
Chen, L., Hasegawa, A.: 1974, Phys. Fluids **17**, 1399
Cheng, C.C., Doschek, G.A., Feldman, U.: 1979, Astrophys. J. **227**, 1037
Chin, Y.C., Wentzel, D.G.: 1972, Astrophys. Space Sci. **16**, 465
Chiuderi, C.: 1981, in Solar Phenomena in Stars and Stellar Systems, R.M. Bonnet, A.K. Dupree Eds., Reidel, Dordrecht, Holland, p. 267
Chiuderi, C.: 1983, in Solar and Magnetic Fields: Origins and Coronal Effects, IAU Symp. **102**, J.O. Stenflo Ed., p. 375
Chiueh, T., Zweibel, E.G.: 1987, Astrophys. J. **317**, 900

Cook, J.W., Brueckner, G.E., Bartoe, J.-D.-F. Socker, D.G.: 1984, Adv. Space Res. **4**, No. 8, 59
Couturier, P., Mangeney, A., Souffrin, P. 1980: in Solar and Interplanetary Dynamics, IAU Symp. **91**, M. Dryer, E. Tandberg-Hanssen Eds., p. 127
Cowling, T.G.: 1953, in The Sun, G.P. Kuiper Ed., Chicago Univ. Press, Chicago, USA, p. 532
Cram, L.E., Damé, L.: 1983, Astrophys. J. **272**, 355
Cram, L.E., Keil, S.L., Ulmschneider, P.: 1979, Astrophys. J. **234**, 768
Cuntz, M., Ulmschneider, P.: 1988, Astron. Astrophys. **193**, 119
Dahlburg, R.B., Dahlburg, J.P., Mariska, J.T.: 1988, Astron. Astrophys. **198**, 300
D'Angelo, N.: 1968, Astrophys. J. **154**, 401
D'Angelo, N.: 1969, Solar Phys. **7**, 321
Davila, J.M.: 1987, Astrophys. J. **317**, 514
Defouw, R.J.: 1976, Astrophys. J. **209**, 266
De Jager, C.: 1959, in Handbuch der Physik **52**, S.F. Flügge Ed., p. 80
De Jager, C., Kuperus, M.: 1961, Bull. Astron. Inst. Neth. **16**, 71
Deming, D.: 1987, in Cool Stars, Stellar Systems and the Sun, J.L. Linsky, R.E. Stencel Eds., Lecture Notes in Physics **291**, Springer, Berlin, Germany, p. 361
Deubner, F.-L.: 1974, Solar Phys. **39**, 31
Deubner, F.-L.: 1975, Astron. Astrophys. **44**, 371
Deubner, F.-L.: 1976, Astron. Astrophys. **51**, 189
Deubner, F.-L.: 1985, in Chromospheric Diagnostics and Modelling, B.W. Lites Ed., Sunspot NM, USA, p. 279
Deubner, F.-L., Durrant, C.J., Kaltenbacher, J.: 1982, Astron. Astrophys. **114**, 85
Deubner, F.-L., Fleck, B.: 1989, Astron. Astrophys., in press
Deubner, F.-L., Reichling, M., Langhanki, R.: 1988, in Advances in Helio- and Asteroseismology, IAU Symp. **123**, J. Christensen-Dalsgaard, S. Frandsen Eds., p. 439
Doschek, G.A., Feldman, U.: 1977, Astrophys. J. **212**, L143
Duijveman, A., Hoyng, P., Ionson, J.A.: 1981, Astrophys. J. **245**, 721
Durrant, C.J.: 1980, Astron. Astrophys. **89**, 80
Durrant, C.J., Michalitsanos, A.G.: 1971, Solar Phys. **18**, 60
Edlén, B.: 1941, Ark. Mat. Astr. Och. Fys. **28B**, No. 1
Einaudi, G., Mok, Y.: 1987, Astrophys. J. **319**, 520
Ferrari, A., Rosner, R., Vaiana, G.S.: 1982, Astrophys. J. **263**, 944
Flå, T., Habbal, S.R., Holzer, T.E., Leer, E.: 1984, Astrophys. J. **280**, 382
Fleck, B., Deubner, F.-L.: 1989, Astron. Astrophys., in press
Flower, D.R.: 1977, Astron. Astrophys. **58**, 185
Flower, D.R., Pineau des Forêts, G.: 1983, Astron. Astrophys. **119**, 321
Fontaine, G., Villeneuve, B., Wilson, J.: 1981, Astron. Astrophys. **243**, 550
Forbes, T.G., Priest, E.R.: 1984, Solar Phys. **94**, 315
Fossat, E., Ricort, G.: 1975, Astron. Astrophys. **43**, 253
Foukal, P.V.: 1976, Astrophys. J. **210**, 575
Foukal, P., Smart, M.: 1981, Solar Phys. **69**, 15
Foukal, P., Hoyt, C., Gilliam, L.: 1986, Astrophys. J. **303**, 861
Galeev, A.A., Rosner, R., Serio, S., Vaiana, G.S.: 1981, Astrophys. J. **243**, 301
Gerassimenko, M., Solodyna, C.V., Nolte, J.T.: 1978, Solar Phys. **57**, 103
Giovanelli, R.G.: 1975, Solar Phys. **44**, 299
Giovanelli, R.G.: 1979, Solar Phys. **62**, 253
Giovanelli, R.G., Beckers, J.M.: 1983, in Solar and Magnetic Fields: Origins and Coronal Effects, IAU Symp. **102**, J.O. Stenflo Ed., p. 407

Giovanelli, R.G., Livingston, W.C., Harvey, J.W.: 1978, Solar Phys. **59**, 49
Glencross, W.M.: 1975, Astrophys. J. **199**, L53
Gokhale, M.H.: 1975, Solar Phys. **41**, 381
Gold, T.: 1964, in The Physics of Solar Flares, W. Hess Ed., NASA SP-50, Greenbelt MD, USA, p. 389
Golub, L.: 1983, in Solar and Magnetic Fields: Origins and Coronal Effects, IAU Symp. **102**, J.O. Stenflo Ed., p. 345
Golub, L. Maxson, C. Rosner, R., Serio, S., Vaiana, G.S.: 1980, Astrophys. J. **238**, 343
Gomez, D., Ferro Fontan, C.: 1988, Solar Phys. **116**, 33
Gonczi, G., Mangeney, A., Souffrin, P.: 1977, Astron. Astrophys. **54**, 689
Gordon, B.E., Hollweg, J.V.: 1983, Astrophys. J. **266**, 373
Grossmann, W., Smith, R.A.: 1988, Astrophys. J. **332**, 476
Grossmann, W., Tataronis, J.: 1973, Z. Phys. **261**, 217
Habbal, S.R., Leer, E.: 1982, Astrophys. J. **253**, 318
Habbal, S.R., Leer, E., Holzer, T.E.: 1979, Solar Phys. **64**, 287
Habbal, S.R., Ronan, R., Withbroe, G.L.: 1985, Solar Phys. **98**, 323
Hammer, R.: 1982, Astrophys. J. **259**, 779
Harrison, R.A.: 1987, Astron. Astrophys. **182**, 337
Harrison, R.A., Pearce, G., Skumanich, A.: 1988, Astrophys. J. **332**, 1058
Hasan, S.S., Schüssler, M.: 1985, Astron. Astrophys. **151**, 69
Hasegawa, A., Chen, L.: 1974, Phys. Rev. Letters **32**, 454
Hasegawa, A., Chen, L.: 1976, Phys. Fluids **19**, 1924
Hasegawa, A., Uberoi, C.: 1982, The Alfvén wave, Tech. Inf. Center, US Dept. of Energy, Oak Ridge, Tenn.
Hasegawa, A.: 1985, in Unstable Current Systems and Plasma Instabilities in Astrophysics, IAU Symp. **107**, M.R. Kundu, G.D. Holman Eds., p. 381
Hearn, A.G., Vardavas, I.M.: 1981, Astron. Astrophys. **98**, 230
Herbold, G., Ulmschneider, P., Spruit, H.C., Rosner, R.: 1985, Astron. Astrophys. **145**, 157
Heyvaerts, J.: 1984, in The Hydromagnetics of the Sun, Proceedings of the Fourth European Meeting on Solar Physics, T.D. Guyenne, J.J. Hunt Eds., ESA SP 220, p. 123
Heyvaerts, J.: 1985, in Unstable Current Systems and Plasma Instabilities in Astrophysics, IAU Symp. **107**, M.R. Kundu, G.D. Holman Eds., p. 95
Heyvaerts, J., Priest, E.R.: 1983, Astron. Astrophys. **117**, 220
Heyvaerts, J., Priest, E.R.: 1984, Astron. Astrophys. **137**, 63
Heyvaerts, J., Priest, E.R., Rust, D.M.: 1977, Astrophys. J. **216**, 123
Hinata, S.: 1980, Astrophys. J. **235**, 258
Hinata, S.: 1981, Astrophys. J. **246**, 532
Hinata, S.: 1983, Solar Phys. **87**, 95
Hinkle, K.H., Hall, D.N.B., Ridgway, S.T.: 1982, Astrophys. J. **252**, 697
Hollweg, J.V.: 1971, Phys. Rev. Letters **27**, 1349
Hollweg, J.V.: 1978a, Solar Phys. **56**, 305
Hollweg, J.V.: 1978b, Geophys. Res. Letters **5**, 731
Hollweg, J.V.: 1981a, in Solar Active Regions, F.Q. Orrall Ed., Colorado Ass. Univ. Press, Boulder CO, USA, p. 277
Hollweg, J.V.: 1981b, Solar Phys. **70**, 25
Hollweg, J.V.: 1982a, Astrophys. J. **254**, 806
Hollweg, J.V.: 1982b, Astrophys. J. **257**, 345
Hollweg, J.V.: 1983, in Solar Wind V, M. Neugebauer Ed., NASA CP-2280, p. 5

Hollweg, J.V.: 1984a, Solar Phys. **91**, 269
Hollweg, J.V.: 1984b, Astrophys. J. **277**, 392
Hollweg, J.V.: 1985a, in Chromospheric Diagnostics and Modelling, B.W. Lites Ed., Sunspot NM, USA, p. 235
Hollweg, J.V.: 1985b, in Advances in Space Plasma Physics, I.C.T.P. Trieste (Italy), B. Buti Ed., World Scientific Publ. Co., Singapore, p. 77
Hollweg, J.V.: 1987a, Astrophys. J. **312**, 880
Hollweg, J.V.: 1987b, Astrophys. J. **317**, 918
Hollweg, J.V.: 1987c, Astrophys. J. **320**, 875
Hollweg, J.V., Jackson, S., Galloway, D.: 1982, Solar Phys. **75**, 35
Hollweg, J.V., Sterling, A.C.: 1984, Astrophys. J. **282**, L31
Hoyle, F.: 1949, Some Recent Researches in Solar Physics, Cambridge Univ. Press, Cambridge, England
Ionson, J.A.: 1978, Astrophys. J. **226**, 650
Ionson, J.A.: 1982, Astrophys. J. **254**, 318
Ionson, J.A.: 1983, Astrophys. J. **271**, 778
Ionson, J.A.: 1984, Astrophys. J. **276**, 357
Ionson, J.A.: 1985a, in Unstable Current Systems and Plasma Instabilities in Astrophysics, IAU Symp. **107**, M.R. Kundu, G.D. Holman Eds., p. 139
Ionson, J.A.: 1985b, Solar Phys. **100**, 289
Ionson, J.A.: 1985c, Astron. Astrophys. **146**, 199
Jeffrey, A.: 1966, Magnetohydrodynamics, Oliver and Boyd, Edinburgh, Scotland
Jones, H.P., Giovanelli, R.G.: 1982, Solar Phys. **79**, 247
Judge, P.G.: 1989, in The Evolution of Peculiar Red Giant Stars, IAU Coll. **106**, in press
Kaburaki, O., Uchida, Y.: 1971, Publ. Astron. Soc. Japan **23**, 405
Kalkofen, W., Ulmschneider, P., Schmitz, F.: 1984, Astrophys. J. **287**, 952
Klein, R.I., Stein, R.F., Kalkofen, W.: 1976, Astrophys. J. **205**, 499
Kohl, J.L., Weiser, H., Withbroe, G.L., Noyes, R.W., Parkinson, W.H., Reeves, E.M., Munro, R.H, Mac Queen, R.M.: 1980, Astrophys. J. **241**, L117
Kopecky, M., Obridko, V.: 1968, Solar Phys. **5**, 354
Kopp, R.A.: 1968, Ph.D. Thesis, Air Force Cambridge Res. Lab., Sci. Rept. **4**, Bedford MA, USA
Koutchmy, S., Zugzda, Y.D., Locans, V.: 1983, Astron. Astrophys. **120**, 185
Kumar, M.: 1987, Ph.D. Thesis, Univ. Meerut, India
Kumar, M., Narain, U.: 1985, Solar Phys. **95**, 69
Kuperus, M.: 1965, Ph.D. Thesis, Univ. Utrecht, Netherlands
Kuperus, M.: 1969, Space Sci. Rev. **9**, 713
Kuperus, M.: 1972, Solar Phys. **22**, 257
Kuperus, M.: 1983, Space Sci. Rev. **34**, 47
Kuperus, M., Ionson, J.A.: 1985, Astron. Astrophys. **148**, 309
Kuperus, M., Ionson, J.A., Spicer, D.S.: 1981, Ann. Rev. Astron. Astrophys. **19**, 7
Lee, M.A.: 1980, Astrophys. J. **240**, 693
Lee, M.A., Roberts, B.: 1986, Astrophys. J. **301**, 430
Leibacher, J.W., Stein, R.F.: 1980, in Highlights of Astronomy **5**, P.A. Wayman Ed., Reidel, Dordrecht, Holland, p. 581
Leibacher, J.W., Stein, R.F.: 1982, in Second Cambridge Workshop on Cool Stars, Stellar Systems and the Sun, M.S. Giampapa, L. Golub Eds., Smithsonian Astrophys. Obs. Special Report No. 392, p. 23
Leighton, R.B., Noyes, R.W., Simon, G.W.: 1962, Astrophys. J. **135**, 474
Leroy, B., Schwartz, S.J.: 1982, Astron. Astrophys. **112**, 84

Levine, R.H.: 1974a, Astrophys. J. **190**, 447
Levine, R.H.: 1974b, Astrophys. J. **190**, 457
Levine, R.H., Withbroe, G.L.: 1977, Solar Phys. **51**, 83
Lighthill, M.J.: 1952, Proc. R. Soc. London **A211**, 564
Lighthill, M.J.: 1954, Proc. R. Soc. London **A222**, 1
Lindsey, C.A.: 1981, Astrophys. J. **244**, 659
Lindsey, C., Kaminski, C.: 1984, Astrophys. J. **282**, L103
Lindsey, C., Roellig, T.: 1987, Astrophys. J. **313**, 877
Linsky, J.L., Ayres, T.R.: 1978, Astrophys. J. **220**, 619
Lites, B.W.: 1984, Astrophys. J. **277**, 874
Lou, Y.Q., Rosner, R.: 1986, Astrophys. J. **309**, 874
Lucy, L.B.: 1982a, Astrophys. J. **255**, 278
Lucy, L.B.: 1982b, Astrophys. J. **255**, 286
Machado, M.E., Moore, R.L.: 1986, Adv. Space Res. **6**, 217
M¨akle, R.: 1969, Solar Phys. **10**, 348
Mariska, J.T., Boris, J.P.: 1983, Astrophys. J. **267**, 409
Mariska, J.T., Hollweg, J.V.: 1985, Astrophys. J. **296**, 746
Martens, P.C.H., Kuin, N.P.M.: 1983, Astron. Astrophys. **123**, 216
Martens, P.C.H., Kuperus, M.: 1982, Astron. Astrophys. **113**, 324
Martens, P.C.H., Van den Oord, G.H.J., Hoyng, P.: 1985, Solar Phys. **96**, 253
McClymont, A.N., Canfield, R.C.: 1983, Astrophys. J. **265**, 483
Mein, N., Schmieder, B.: 1981, Astron. Astrophys. **97**, 310
Mein, P.: 1981, in Solar Phenomena in Stars and Stellar Systems, R.M. Bonnet, A.K. Dupree Eds., Reidel, Dordrecht, Holland, p. 265
Merzljakov, E.G., Ruderman, M.S.: 1986, Solar Phys. **103**, 259
Mewe, R.: 1979, Space Sci. Rev. **24**, 101
Middelkoop, F.: 1982, Astron. Astrophys. **113**, 1
Milkey, R.W.: 1970, Solar Phys. **14**, 62
Mok, Y.: 1987, Astron. Astrophys. **172**, 327
Mok, Y., Einaudi, G.: 1985, J. Plasma Phys. **33**, 199
Mullan, D.J.: 1984, Astrophys. J. **282**, 603
Musielak, Z.E.: 1982, Astron. Astrophys. **105**, 23
Musielak, Z.E.: 1987, Astrophys. J. **322**, 234
Musielak, Z.E., Rosner, R.: 1987, Astrophys. J. **315**, 371
Musielak, Z.E., Rosner, R.: 1988, Astrophys. J. **329**, 376
Musielak, Z.E., Rosner, R., Ulmschneider, P.: 1987, in Cool Stars, Stellar Systems and the Sun, J.L. Linsky, R.E. Stencel Eds., Lecture Notes in Physics **291**, Springer, Berlin, Germany, p. 66
Musielak, Z.E., Rosner, R., Ulmschneider, P.: 1989, Astrophys. J. in press
Narain, U., Kumar, M.: 1985, Solar Phys. **99**, 111
Narain, U., Kumar, M.: 1986, Astron. Astrophys. **164**, 373
Nicolas, K.R., Bartoe, J.-D.-F., Brueckner, G.E., Van Hoosier, M.E. 1979, Astrophys. J. **233**, 741
Nocera, L., Leroy, B., Priest, E.R.: 1984, Astron. Astrophys. **133**, 387
Nocera, L., Leroy, B., Priest, E.R.: 1985, in Unstable Current Systems and Plasma Instabilities in Astrophysics, IAU Symp. **107**, M.R. Kundu, G.D. Holman Eds., p. 365
Nocera, L., Priest, E.R.: 1984, in The Hydromagnetics of the Sun, Proceedings of the Fourth European Meeting on Solar Physics, T.D. Guyenne, J.J. Hunt Eds., ESA SP-220, p. 249
Nolte, J.T., Gerassimenko, M., Krieger, A.S., Petrasso, R.D., Svestka, Z., Wentzel, D.G.:

1977, Solar Phys. **55**, 401

Nowak, T., Ulmschneider, P.: 1977, Astron. Astrophys. **60**, 413

Oranje, B.J., Zwaan, C.: 1985, Astron. Astrophys. **147**, 265

Osterbrock, D.E.: 1961, Astrophys. J. **134**, 347

Owocki, S.P., Castor, J.I., Rybicki, G.B.: 1988, Astrophys. J., in press

Pallavicini, R., Golub, L., Rosner, R., Vaiana, G.S., Ayres, T., Linsky J.L.: 1981, Astrophys. J. **248**, 279

Parker, E.N.: 1972, Astrophys. J. **174**, 499

Parker, E.N.: 1975, Astrophys. J. **201**, 494

Parker, E.N.: 1979, Astrophys. J. **230**, 905

Parker, E.N.: 1981a, Astrophys. J. **244**, 631

Parker, E.N.: 1981b, Astrophys. J. **244**, 644

Parker, E.N.: 1981c, in Solar Phenomena in Stars and Stellar Systems, R.M. Bonnet, A.K. Dupree Eds., Reidel, Dordrecht, Holland, p. 33

Parker, E.N.: 1982, Geophys. Astrophys. Fluid Dyn. **22**, 195

Parker, E.N.: 1983a, Astrophys. J. **264**, 635

Parker, E.N.: 1983b, Astrophys. J. **264**, 642

Parker, E.N.: 1983c, in Solar Wind V, M. Neugebauer Ed., NASA CP-2280, p. 23

Parker, E.N.: 1985, Geophys. Astrophys. Fluid Dyn. **34**, 243

Parker, E.N.: 1986, in Coronal and Prominence Plasmas, A.I. Poland Ed., NASA CP-2442, p. 9

Parker, E.N.: 1988, Astrophys. J. **330**, 474

Pasachoff, J.M., Ladd, E.F.: 1987, Solar Phys. **109**, 365

Pasachoff, J.M., Landman, D.A.: 1984, Solar Phys. **90**, 325

Piddington, J.H.: 1956, Mon. Not. Roy. Astron. Soc., **116**, 314

Piddington, J.H.: 1973, Solar Phys. **33**, 363

Pineau des Forêts, G.: 1979, Astron. Astrophys. **78**, 159

Poedts, S., Goossens, M.: 1987, Solar Phys. **109**, 265

Poedts, S., Goossens, M.: 1988, Astron. Astrophys. **198**, 331

Porter, J.G., Moore, R.L.: 1987, in Proceedings of the 9th Sacramento Peak Summer Workshop "Solar and Stellar Coronal Structure and Dynamics", Sunspot NM, USA

Porter, J.G., Moore, R.L., Reichmann, E.J., Engvold, O., Harvey, K.L. 1987, Astrophys. J. **323**, 380

Porter, J.G., Toomre, J., Gebbie, K.B.: 1984, Astrophys. J **283**, 879

Priest, E.R.: 1981, in Solar Flare Magnetohydrodynamics, E.R. Priest Ed., Gordon and Breach, New York, p. 139

Priest, E.R.: 1982, Solar Magnetohydrodynamics, Reidel, Dordrecht, Holland

Priest, E.R.: 1983, Solar Phys. **86**, 33

Priest, E.R., Raadu, M.A.: 1975, Solar Phys. **43**, 177

Provost, J., Mein, N.: 1979, Solar Phys. **64**, 43

Querci, M., Querci, F.: 1985, Astron. Astrophys. **147**, 121

Rabin, D., Moore, R.: 1984, Astrophys. J. **285**, 359

Rae, I.C., Roberts, B.: 1982, Mon. Not. Roy. Astron. Soc. **201**, 1171

Renzini, A., Cacciari, C., Ulmschneider, P., Schmitz, F.: 1977, Astron. Astrophys. **61**, 39

Revathy, P.: 1977, Solar Phys. **53**, 445

Roberts, B.: 1983, in Solar and Magnetic Fields: Origins and Coronal Effects, IAU Symp. **102**, J.O. Stenflo Ed., p. 61

Roberts, B., Mangeney, A.: 1982, Mon. Not. Roy. Astr. Soc. **198**, 7P

Roberts, B., Webb, A.R.: 1978, Solar Phys. **56**, 5

Rosner, R., Golub, L. Coppi, B., Vaiana, G.S.: 1978a, Astrophys. J. **222**, 317

Rosner, R., Low, B.C., Holzer, T.E.: 1986, in Physics of the Sun, Vol. II, P.A. Sturrock Ed., Reidel, Dordrecht, Holland, p. 135
Rosner, R., Tucker, W.H., Vaiana, G.S.: 1978b, Astrophys. J. **220**, 643
Rutten, R.G.M.: 1987, Astron. Astrophys. **177**, 131
Rybicki, G.B.: 1987, in Instabilities in Luminous Early Type Stars, H.J.G.L.M. Lamers, C.W.H. de Loore Eds., Reidel, Dordrecht, Holland, p.175
Saar, S.H.: 1987, in Cool Stars, Stellar Systems and the Sun, J.L. Linsky, R.E. Stencel Eds., Lecture Notes in Physics **291**, Springer, Berlin, Germany, p. 10
Sagdeev, R.Z., Galeev, A.A.: 1969, Nonlinear Plasma Theory, Benjamin, New York, USA
Sakurai, T. Granik, A.: 1984, Astrophys. J. **277**, 404
Sakurai, T., Levine, R.H.: 1981, Astrophys. J. **248**, 817
Schatten, K.H., Mayr, H.G.: 1986, Astrophys. J. **309**, 864
Schatzman, E.: 1949, Ann. d'Astrophys. **12**, 203
Schatzman, E., Souffrin, P.: 1967, Ann. Rev. Astron. Astrophys. **5**, 67
Scheurwater, R., Kuperus, M.: 1988, Astron. Astrophys. **194**, 213
Schmieder, B.: 1979, Astron. Astrophys. **74**, 273
Schmieder, B., Mein, N.: 1980, Astron. Astrophys. **84**, 99
Schmitz, F., Ulmschneider, P.: 1980a, Astron. Astrophys. **84**, 93
Schmitz, F., Ulmschneider, P.: 1980b, Astron. Astrophys. **84**, 191
Schmitz, F., Ulmschneider, P.: 1981, Astron. Astrophys. **93**, 178
Schmitz, F., Ulmschneider, P., Kalkofen, W.: 1985, Astron. Astrophys. **148**, 217
Schrijver, C.J.: 1987a, Astron. Astrophys. **172**, 111
Schrijver, C.J.: 1987b, Astron. Astrophys. **180**, 241
Schrijver, C.J.: 1987c, in Cool Stars, Stellar Systems and the Sun, J.L. Linsky, R.E. Stencel Eds., Lecture Notes in Physics **291**, Springer, Berlin, Germany, p. 135
Schrijver, C.J., Coté, J., Zwaan, C., Saar, S.H.: 1989, Astron. Astrophys. **337**, 964
Schrijver, C.J., Mewe, R, Walter, F.M.: 1984, Astron. Astrophys. **138**, 258
Schüssler, M.: 1984, Astron. Astrophys. **140**, 453
Schwartz, S.J., Bel, N.: 1984, Astron. Astrophys. **137**, 128
Schwartz, S.J., Cally, P.S., Bel, N.: 1984, Solar Phys. **92**, 81
Schwartz, S.J., Leroy, B.: 1982, Astron. Astrophys. **112**, 93
Schwarzschild, M.: 1948, Astrophys. J. **107**, 1
Shore, S.N., Brown, D.N., Sonneborn, G., Gibson, D.M.: 1987, Astron. Astrophys. **182**, 285
Similon, P.L., Sudan, R.N.: 1989, Astrophys. J. **336**, 442
Skumanich, A.: 1972, Astrophys. J. **171**, 565
Somasundaram, K., Uberoi, C.: 1982, Solar Phys. **81**, 19
Somov, B.V., Syrovatskii, S.I.: 1977, Solar Phys. **55**, 393
Souffrin, P.: 1982, Astron. Astrophys. **109**, 205
Spicer, D.S.: 1982, Space Sci. Rev. **31**, 351
Spicer, D.S.: 1983, Solar Phys. **88**, 43
Spruit, H.C.: 1981, Astron. Astrophys. **98**, 155
Spruit, H.C.: 1982, Solar Phys. **75**, 3
Stefanik, R.: 1969, Ph.D. Thesis, Air Force Cambridge Res. Lab., Sci. Rept. **5**, Bedford MA, USA
Stéfant, R.J.: 1970, Phys. Fluids **13**, 440
Stein, R.F.: 1967, Solar Phys. **2**, 385
Stein, R.F.: 1968, Astrophys. J. **154**, 297
Stein, R.F.: 1969, in Chromosphere-Corona Transition Region, R.G. Athay, G. Newkirk Eds., High Altitude Observatory, Boulder CO, USA, p. 171

Stein, R.F.: 1981, Astrophys. J. **246**, 966
Stein, R.F., Leibacher, J.: 1974, Ann. Rev. Astron. Astrophys. **12**, 407
Stein, R.F., Leibacher, J.: 1980, in Stellar Turbulence, D.F. Gray, J.L. Linsky Eds., IAU Colloq. No. 51, Lecture Notes in Physics **114**, Springer, Berlin, Germany, p. 225
Stein, R.F., Schwartz, R.A.: 1972, Astrophys. J. **177**, 807
Steinolfson, R.S., Priest, E.R., Poedts, S., Nocera, L., Goossens, M.: 1986, Astrophys. J. **304**, 526
Steinolfson, R.S., Van Hoven, G.: 1984, Astrophys. J. **276**, 391
Stencel, R.E., Mullan, D.J., Linsky, J.L., Basri, G.S., Worden, S.P. 1980, Astrophys. J. Suppl. **44**, 383
Stenflo, J.O.: 1978, Rep. Prog. Phys. **75**, 3
Sterling, A.C., Hollweg, J.V.: 1988, Astrophys. J. **327**, 950
Strauss, H.R.: 1988, Astrophys. J. **326**, 412
Strauss, H.R., Otani, N.F.: 1988, Astrophys. J. **326**, 418
Sturrock, P.A., Uchida, Y.: 1981, Astrophys. J. **246**, 331
Syrovatskii, S.I.: 1978, Solar Phys. **58**, 89
Taylor, J.B.: 1974, Phys. Rev. Letters **33**, 1139
Tucker, W.H.: 1973, Astrophys. J. **186**, 285
Tur, T.J., Priest, E.R.: 1976, Solar Phys. **48**, 89
Uberoi, C.: 1982, Solar Phys. **78**, 351
Uchida, Y.: 1963, Publ. Astron. Soc. Japan **15**, 376
Uchida, Y., Kaburaki, O.: 1974, Solar Phys. **35**, 451
Uchida, Y., Sakurai, T.: 1977, Solar Phys. **51**, 413
Ulmschneider, P.: 1967, Z. Astrophys. **67**, 193
Ulmschneider, P.: 1970, Solar Phys. **12**, 403
Ulmschneider, P.: 1971a, Astron. Astrophys. **12**, 297
Ulmschneider, P.: 1971b, Astron. Astrophys. **14**, 275
Ulmschneider, P.: 1974, Solar Phys. **39**, 327
Ulmschneider, P.: 1976, Solar Phys. **49**, 249
Ulmschneider, P.: 1979, Space Sci. Rev. **24**, 71
Ulmschneider, P.: 1981, in Solar Phenomena in Stars and Stellar Systems, R.M. Bonnet, A.K. Dupree Eds., Reidel, Dordrecht, Holland, p. 239
Ulmschneider, P.: 1986, Adv. Space Res. **6**, No. 8, 39
Ulmschneider, P.: 1988, Astron. Astrophys. **197**, 223
Ulmschneider, P.: 1989, Astron. Astrophys. in press
Ulmschneider, P., Muchmore, D.: 1986, in Small Scale Magnetic Flux Concentrations in the Solar Photosphere, W. Deinzer, M. Knölker, H.H. Voigt Eds., Vandenhoeck and Ruprecht, Göttingen, Germany, p. 191
Ulmschneider, P., Muchmore, D., Kalkofen, W.: 1987, Astron. Astrophys. **177**, 292
Ulmschneider, P., Schmitz, F., Hammer, R.: 1979, Astron. Astrophys. **74**, 229
Ulmschneider, P., Schmitz, F., Kalkofen, W., Bohn, H.U.: 1978, Astron. Astrophys. **70**, 487
Ulmschneider, P., Stein, R.F.: 1982, Astron. Astrophys. **106**, 9
Ulmschneider, P., Zähringer, K.: 1987, in Cool Stars, Stellar Systems and the Sun, J.L. Linsky, R.E. Stencel Eds., Lecture Notes in Physics **291**, Springer, Berlin, Germany, p. 63
Vaiana, G.S., et al.: 1981, Astrophys. J. **245**, 163
Vainshtein, S.I., Parker, E.N.: 1986, Astrophys. J. **304**, 821
Van Ballegooijen, A.A.: 1985, Astrophys. J. **298**, 421
Van Ballegooijen, A.A.: 1986, Astrophys. J. **311**, 1001

Van de Hulst, H.C.: 1953, in The Sun, G.P. Kuiper Ed., Chicago Univ. Press, Chicago, USA, p. 207
Vekstein, G.E.: 1987, Astron. Astrophys. **182**, 324
Vernazza, J.E., Avrett, E.H., Loeser, R.: 1981, Astrophys. J. Suppl. **45**, 635
Vernazza, J.E., Foukal, P.V., Huber, M.C.E., Noyes, R.W., Reeves, E.M., Schmahl, E.J., Timothy, J.G., Withbroe, G.L. 1975, Astrophys. J. **199**, L123
Vilhu, O.: 1987, in Cool Stars, Stellar Systems and the Sun, J.L. Linsky, R.E. Stencel Eds., Lecture Notes in Physics **291**, Springer, Berlin, Germany, p. 110
von Uexküll, M., Kneer, F., Malherbe, J.M., Mein, P.: 1989, Astron. Astrophys. **208**, 290
Walter, F.M., Schrijver, C.J.: 1987, in Cool Stars, Stellar Systems and the Sun, J.L. Linsky, R.E. Stencel Eds., Lecture Notes in Physics **291**, Springer, Berlin, Germany, p. 262
Wentzel, D.G.: 1974, Solar Phys. **39**, 129
Wentzel, D.G.: 1976, Solar Phys. **50**, 343
Wentzel, D.G.: 1977, Solar Phys. **52**, 163
Wentzel, D.G.: 1978, Rev. Geophys. Space Phys. **16**, 757
Wentzel, D.G.: 1979a, Astron. Astrophys. **76**, 20
Wentzel, D.G.: 1979b, Astrophys. J. **227**, 319
Wentzel, D.G.: 1979c, Astrophys. J. **233**, 756
Wentzel, D.G.: 1981, in The Sun as a Star, Monograph Series on Nonthermal Phenomena in Stellar Atmospheres, S. Jordan Ed., NASA SP-450, p. 331
Wentzel, D.G.: 1989, Astrophys. J. **336**, 1073
Whitehouse, D.R.: 1985, Astron. Astrophys. **145**, 449
Wilson, O.C.: 1963, Astrophys. J. **138**, 832
Wilson, O.C., Bappu, M.K.V.: 1957, Astrophys. J. **125**, 661
Withbroe, G.L., Noyes, R.W.: 1977, Ann. Rev. Astron. Astrophys. **15**, 363
Withbroe, G.L., Kohl, J.L., Weiser, H., Munro, R.H.: 1985, Astrophys. J. **297**, 324
Withbroe, G.L.: 1988, Astrophys. J. **325**, 442
Wolf, B.E.: 1987, Astron. Astrophys. **179**, 371
Woltjer, L.: 1958, Bull. Astron. Inst. Neth. **14**, 39
Zähringer, K., Ulmschneider, P.: 1987, in The Role of Fine- Scale Magnetic Fields on the Structure of the Solar Atmosphere, E.H. Schröter, M. Vazquez, A.A. Wyller Eds., Cambridge Univ. Press, Cambridge, England, p. 243
Zuccarello, F., Burm, H., Kuperus, M., Raadu, M., Spicer, D.S.: 1987, Astron. Astrophys. **180**, 218
Zugzda, Y.D., Locans, V.: 1982, Solar Phys. **76**, 77
Zwaan, C.: 1978, Solar. Phys. **60**, 213
Zweibel, E.: 1980, Solar Phys. **66**, 305