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ON THE THEORY OF CORONAL HEATING MECHANISMS

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## I. INTRODUCTION

The theory of the formation of the solar chromosphere and corona has as its main challenging goal to specify the nature of the heating mechanism. Since the discovery of the ( $\sim 10^6$ K) hot solar corona about forty years ago the theory of coronal heating has received vivid interest. Early theories were essentially based on the dissipation of acoustic waves or weak shock waves generated by the photospheric granulation (Biermann, 1946; 1948; Schwarzschild, 1948; Schatzman, 1949). They looked very much like a description of the extension of the photospheric model atmosphere by introducing in addition to the radiative energy flux and the opacity a mechanical energy flux and a wave absorption coefficient. The atmosphere was assumed to be largely spherically symmetric and all the model calculations were one dimensional and steady, specifying boundary conditions at appropriate levels in the atmosphere.

It was soon realized that the corona overlying an active region needed to be heated one to two orders of magnitude more efficiently than the so-called quiet corona and because of the active region's magnetic nature it became evident that magnetic fields must play an important role in the energetics.

Another important discovery was the solar wind which is a direct consequence of the existence of a hot corona and thus of coronal heating. Spaceborne observations of the solar wind made it clear that actual irreversible heating as well as momentum addition must be present up to large distances from the sun, far beyond the level of the maximum coronal temperature.

It is unlikely that weak shock waves can reach that far. They are probably exhausted already in the inner corona. Here the magnetic field must play a dominant role in selecting the required mode of energy transfer, presumably

ly in the form of Alfvén waves.

Occasionally the corona is heated to very high temperatures ( $10^7 - 10^8$  K) during a solar flare. It is generally accepted that Joule heating and magnetic reconnection, conversion of stored magnetic energy into heat and high energetic particles, is the fundamental flare process, and wave dissipation plays a negligible role in this case.

We are thus facing two basically different mechanisms of coronal heating:

- a. Heating by acoustic processes in the "non magnetic" parts of the atmosphere ( $\beta = (8\pi p/B^2) \gtrsim 1$ ) of which the shock wave theory is an example worked out in great deal.
- b. Heating by electrodynamic processes in the magnetic regions of the corona. ( $\beta \ll 1$ ) either by
  - 1) magnetohydrodynamic waves or
  - 2) current heating in the regions with large electric current densities (flare type heating) .

The present state of the art in these two classes of theories and the typical problems one encounters in each of them will be discussed in sections 3, 4.

In a sense the shock wave theory was rather successful in that it linked the observed photospheric motions to the observed and predicted chromospheric and coronal energy losses known at that time. However, the amount of mechanical energy generated in the sub-photospheric layers is known with an accuracy less than an order of magnitude and so are the radiative losses of the chromosphere and the corona.

For an extensive discussion of models based on shock wave theory the

reader is referred to a great number of (revised) papers (Kuperus, 1969; Bray and Loughhead, 1974; Kaplan et al., 1974; Athay, 1975; Kuperus and Ghislerri, 1977; Withbroe and Noyes, 1977; Flower and Pineau des Forêts, 1977; Ulmschneider, 1967, 1970, 1971).

At present the theory of the formation of the outer solar atmosphere should take into account the following reasonably well established facts.

1. The corona consists of essentially two types of radiation structures which could be associated with or coincident with magnetic structures. (Vaiana and Rosner, 1978).

a) OPEN STRUCTURES, the so called coronal holes, which are thought to be the seats of high speed solar wind streams and have a very low emission measure in soft X-rays (Zirker, 1977).

b) CLOSED STRUCTURES (loops) that radiate more energy than the surrounding corona.

The loops are strongly concentrated in the active regions but also interconnecting loops are present between several centres of activity.

2. The chromospheric network and the plage regions are the seats of strong magnetic fields ( $\sim 2000$  Gauss). The field appears to be concentrated in flux tubes in the lower parts of the atmosphere. They appear to have the same structure underneath the coronal holes as elsewhere (Stenflo, 1978; Zwaan, 1978; Spruit, 1980).

3. The existence of a thin transition region between the chromosphere and the corona with steep temperature gradients and large "turbulent" velocities confirmed by the XUV observations (Withbroe and Noyes,

1977; Brueckner, 1980).

4. A highly filamentary chromosphere, with upward motions in the spicules and downflows near the network boundaries (Brueckner, 1980).
5. A photospheric/chromospheric velocity field with turbulent properties which apparently are evanescent acoustic waves that have a spectrum peaked around a period of 300 sec. (Deubner, 1977; Stein and Leihbacher, 1974).
6. The detection of a strongly varying multicomponent solar wind plasma at the Earth's orbit with large differences between the electron and proton temperature, a very distinct magnetic sector boundary structure and some evidence of the presence of Alfvén waves (Belcher and Davies, 1971).

In the next section the requirements of chromospheric and coronal heating will be discussed in the context of the fundamental constraints one encounters in modelling the outer solar atmosphere.

## II. ENERGY TRANSFER IN THE CHROMOSPHERE AND CORONA

### II.1. Fundamental Equations and Energy Balance

Modelling the outer solar atmosphere involves solving the equations of continuity, momentum and energy in a given geometry usually specified by the magnetic field  $\underline{B}$  and the acceleration of gravity  $\underline{g}$ . The continuity equation is  $\frac{\partial \rho}{\partial t} = -\nabla \cdot \rho \underline{v}$ , where  $\rho$  is the mass density and  $\underline{v}$  the velocity. The equation of motion is

$$\rho \frac{\partial \underline{v}}{\partial t} + \rho (\underline{v} \cdot \nabla) \underline{v} = -\nabla p + \rho \underline{g} + \frac{1}{c} (\underline{j} \times \underline{B}) + \nabla \cdot \underline{\tau}, \quad (1)$$

where  $p = RT/\mu$  is the gas pressure,  $T$  is temperature,  $R$  is the gas constant,  $\mu$  the mean molecular weight.  $\underline{g} = -\nabla \phi$ , where  $\phi$  is the gravitational potential,  $\underline{B}$  the magnetic field vector,  $\underline{j}$  is the current density and  $\underline{\tau}$  is the viscous stress tensor (Landau and Lifshitz 1959).

The equation for the thermodynamic energy  $\underline{U} = \frac{1}{2} \rho \underline{v}^2 + \rho e + \rho \phi$ , where  $e = p / (\gamma - 1)$  is the internal energy, is

$$\frac{\partial \underline{U}}{\partial t} = -\nabla \cdot \underline{F} + \underline{j} \cdot \underline{E}, \quad (2)$$

where the total thermodynamic energy flux  $\underline{F}$  is given by

$$\underline{F} = \rho \underline{v} \left( \frac{1}{2} \underline{v}^2 + e + \phi \right) + p \underline{v} + \underline{\tau} \cdot \underline{v} - \kappa \nabla T + \underline{F}_{\text{RAD}} \quad (3)$$

The terms on the right hand side are respectively the convective energy flux  $\underline{F}_{\text{conv}}$ , the mechanical energy flux  $\underline{F}_m$ , the viscous energy flux  $\underline{F}_{\text{visc}}$ , the conductive energy flux  $\underline{F}_{\text{cond}}$ , and the radiative energy flux  $\underline{F}_{\text{rad}}$ .  $\kappa$  is the coefficient of the thermal conduction, which will be discussed in the next subsection.

The term  $\underline{j} \cdot \underline{E}$  is the electrodynamic energy input representing the

coupling between the thermodynamic system and the electrodynamic system.  $\underline{E} = \eta \underline{j} - (\underline{v} \times \underline{B}) c^{-1}$  is the total electric field and  $\eta$  the electrical resistivity. Instead of writing the energy equation for the thermodynamic energy one could write the equation for the total energy  $U_{\text{tot}}$  including the magnetic energy. Using Poynting's theorem the energy equation can be written in conservative form  $\partial U_{\text{tot}} / \partial t = -\nabla \cdot \underline{F}$ , by adding the Poynting flux  $(\underline{E} \times \underline{B}) c / 4\pi$  to the total flux. We prefer to write the energy equation in the form of equation (2) clearly separating the electrodynamic coupling term  $\underline{j} \cdot \underline{E}$ .

Viscous coupling in the outer layers has not been considered since small wavelengths are required for the viscosity to become important. The viscous flux will be neglected further. In a non magnetic atmosphere in steady state the energy equation (2) is thus reduced to  $\nabla \cdot (\underline{F}_{\text{conv}} + \underline{F}_{\text{mech}} + \underline{F}_{\text{cond}} + \underline{F}_{\text{rad}}) = 0$ .

The energy requirements for the heating of the chromosphere and corona are based on the strongly model dependent estimates of the energy losses in radiation, conduction and convection. Extensive discussions of this problem are found in Athay (1976) and Withbroe and Noyes (1977). We here confine ourselves with summarizing the total

estimated energy losses of the chromosphere and corona in the quiet sun, the coronal hole and the active region as has been given by Withbroe and Noyes (1977).

These energy fluxes have to be replenished by acoustical and/or electrodynamic heating. A flux in between  $10^5 - 10^7$  ergs /  $\text{cm}^2$  sec is necessary depending on the degree of activity and the magnetic topology in the corona. The energy demands are the smallest for the quiet sun although the radiative losses are an order of magnitude larger than in coronal holes. The large energy flux required for a coronal hole essentially results from the large convective losses out of the presumably open magnetic structures. Both the quiet sun and active region consist of closed structures and thus the losses are primarily radiative and conductive.

## II.2. *Microscopic Effects and Transport Parameters*

We cannot over emphasize the care one needs to exercise when one considers the energy and momentum balance problem in a magnetized plasma. This arises for four reasons: (i) the magnetic field can result in highly anisotropic transport coefficients; (ii) some of the ED heating mechanisms cause localized heating which can be isolated by the magnetic field surrounding the heating locus; (iii) some of the ED heating mechanisms can lead to substantial modifications of the transport coefficients themselves - at least locally - and these modifications must be accounted for self-consistently when treating the energy and momentum balance question; and (iv) some ED mechanisms are long wavelength mechanisms (a mechanism that dissipates energy over a substantial fraction of the magnetic configuration parallel to  $\underline{B}$ ) while some ED mechanisms are short wavelength mechanisms (a mechanism that is highly localized parallel to  $\underline{B}$ ) (Spicer and Brown 1980). Such

behaviour on the part of the heating mechanisms can significantly alter the energy balance arguments and resulting scaling laws (Spicer et al 1980).

If the heating mechanism operates along the magnetic field over a length  $L_{\parallel} \gg L_{\perp}$ , which is the characteristic length across the magnetic field, then the ratio of the divergence of the heat flux parallel and perpendicular to the magnetic field can be of the order one. Such a situation applies to the tearing mode process (Spicer 1980) (IV). In the case of anomalous Joule heating this ratio may be much smaller than one.

The above arguments should be contrasted with the classical treatment of thermal conduction (Chapman 1954; Braginskii 1965; Kopp 1968). There it is shown that the heat flux parallel to  $\underline{B}$  is given by

$$\underline{F}_{\text{cond},\parallel} = -\kappa_0 T^{5/2} \nabla_{\parallel} T, \text{ where } \kappa_0 T^{5/2} = 6 \times 10^{-7} T^{5/2} \text{ erg cm}^{-1} \text{ sec}^{-1} \text{ K}^{-1}.$$

This flux is primarily directed along  $\underline{B}$ , unless  $L_{\perp}/L_{\parallel} \ll 1$ .

This was first pointed out and applied to magnetic structures in the chromosphere and corona by Kopp and Kuperus (1968) and later reexamined by Gabriel (1976).

The classical approximation neglects the collective fluctuations. However, when a non-thermal feature imposes itself on a plasma the collective fluctuations can be driven to higher amplitudes. These collective fluctuations can then produce large stochastic particle deflections which thus enhance the effective collision frequency thus leading to anomalous transport parameters. The non-thermal feature that is of present interest is the drift velocity of the current (IV) which must reach a magnitude comparable to the phase velocity of a normal mode of the plasma.

Another effect that is microscopic in origin but also manifests itself macroscopically is the divergence of the radiation flux  $\nabla \cdot \underline{F}_{\text{RAD}}$ , which can be both a sink and a source term. This point is generally

overlooked in energy balance arguments (e.g., Withbroe and Noyes 1977) even though efflorescent heating of the chromosphere by  $\sim 1$  keV soft X-rays emitted in the corona by flares is known to significantly alter the energy balance arguments in solar flares (Henoux and Nakagawa 1977; Henoux and Rust 1980). Similar effects could also occur during quiescent periods and could radically alter the present energy balance arguments (e.g., Withbroe and Noyes 1977).

However, since the corona is to a large degree optically thin one may use the radiative loss function  $\nabla \cdot \underline{F}_{\text{RAD}} = \rho^n f(T)$ , where  $f(T)$  has been calculated by Cox and Tucker (1969) and later refined by McWhirter et al (1975).

An additional microscopic difficulty associated with some ED mechanisms is the problem of how energy released in a localized volume  $\Delta V$  is redistributed throughout the total volume  $V$  of the magnetic configuration such that  $\Delta V/V \ll 1$ . Physically the problem is to determine which transport mechanism is best suited to transport energy perpendicular to  $B$  from the locus of energy release into the remaining volume  $V$ . Is it diffusion or convection? Diffusion processes are rather unlikely due to the large timescales compared to the characteristic cooling timescales. This latter point is true even with convective transport. However, convective transport which leads to much shorter timescales can accomplish this by a variety of means usually depending on the ED heating mechanism. For example, tearing modes (IV) could lead to magnetic braiding as has been discussed in connection with flares (Spicer 1976, 1977) and thereby lead to a substantial increase in the heated volume over the actual energy release volume  $\Delta V$ . On the other hand local energy release must lead to waves being generated at the site of energy release. These waves can then propagate

and damp throughout the volume  $V$  thus heating the plasma as they proceed. This approach has been examined by Spicer et al (1980).

### II.3. *Global Energy Balance and Macroscopic Effects of the Magnetic Field*

II.3.1. GLOBAL ENERGY BALANCE AND NON STATIONARY HEATING The formation of coronal structures requires a detailed knowledge of the local heating throughout the whole structure. Nonetheless several global calculations have been presented balancing the energy input against the total losses in order to find the dependence of the coronal temperature on the energy input.

Rosner et al (1978a) derived a scaling law for coronal loops:

$T \sim (pL)^{1/3}$ , where  $L$  is the loop length and  $p$  the pressure in the loop. Hearn (1975, 1979) suggested that at a given base pressure and a given temperature the corona will adjust in such a way that it picks out the solution with a minimum energy flux, thus reducing the number of parameters.

However, Endler et al (1979) demonstrate that it is the details of the heating mechanism expressed for instance in the dissipation length that determine the ultimate coronal configuration. For instance the height where the maximum temperature is reached strongly depends on the dissipation length (Ulmschneider 1971; McWhirter et al 1975). The basic problem with the minimum flux assumption is that it seems to be over-determined. Once  $p_0$  has been selected the radiative losses are essentially determined and so is the energy input. Changing the energy input will result in a shift of the transition layer to a pressure such that energy balance is maintained (Martens 1980; Van Tend 1980). How this occurs can only be found after solving the time dependent equations.

Even in the simple case of plane geometry such has not yet been done in any satisfactory way. Kuperus and Athay (1967) argued that a very steep lower transition layer is very hard to keep in hydrostatic equilibrium. The argument is that the high coronal temperature is a natural consequence of the inability of the upper chromosphere to radiate the dissipated energy in a low temperature configuration. If it were not for the thermal conduction two separated but adjacent atmospheric layers could exist both in hydrostatic equilibrium assuming pressure continuity. However, the coefficient of heat conduction very abruptly decreases going inwards from the transition layer to the chromosphere. Therefore, if thermal conduction is the essential energetic link between the corona and the upper chromosphere, the upper chromosphere, which is supposed to radiate at its maximum efficiency is likely to be excessively heated. As a consequence the upper layers may be set into motion. Kuperus and Athay thought that this overheating in the upper chromosphere would be the origin of the spicules. It was demonstrated by Kopp and Kuperus (1968) that overheating aggravates in converging magnetic fields due to the magnetic channelling. Converging flux also increases the likeliness of heat flux instability (Spicer 1979). This would then argue in favour of seeing the spicules concentrated on the boundaries where one finds the magnetic knots.

It is tempting to assign the observed large non thermal motions in the transition layer (Brueckner 1980) to the effect of dynamic adjustment of the transition layer to time dependant coronal heating.

On the other hand Pneuman and Kopp (1980) showed that the observed downflows in the transition region carry an enthalpyflux larger than the conductive flux so that in those regions the models based on conductive energy supply need to be reevaluated.

Static modelling is likely to be of little value and might even lead to erroneous conclusions.

Recently the magnitude of the temperature gradient in the transition zone has been questioned on theoretical grounds (Spicer 1979) and on empirical grounds (Nicolas et al 1979).

Using empirical results Spicer (1979) has argued that because the mean free path of the electrons in the transition zone is comparable to the empirically determined characteristic temperature gradient scale length there the heat flux will be marginally unstable electrostatically if the empirical models are correct. He showed that the heat flux would be reduced by ~ 30% over its classical value. Spicer further argued that this reduced heat flux should appear as convective turbulence in the transition zone consistent with the previous arguments of Kuperus and Athay (1967).

On the other hand Nicolas et al (1979) using high resolution data (~ 1") from the HRTS instrument (Bartoe and Brueckner 1975) have brought into question the assumption of taking the filling factor equal to 1 for slit instruments. They conclude by suggesting that most of the transition zone emission may originate from structures whose filling factor would be less than one, hence reducing the temperature to be expected.

II.3.2. MACROSCOPIC EFFECTS OF THE MAGNETIC FIELD As already discussed the magnetic field  $\tilde{B}$  plays a crucial role in the coronal heating process. However, it is not yet clear whether this role is passive, that is  $\tilde{B}$  simply acts as a guide for the transport of energy and momentum, or whether this role is active, that is  $\tilde{B}$  itself is intimately involved in the heating process. This delineation of roles for  $\tilde{B}$  naturally separates the two classes of coronal heating mechanisms:

acoustic (III) and electrodynamic (IV). However, such a precise delineation of roles for  $\underline{B}$  is quite artificial since for  $\underline{B}$  not to play a role in the mechanical class of coronal heating mechanisms requires the acoustic waves involved to propagate exactly parallel to  $\underline{B}$ , since oblique propagation will lead to magnetic field perturbations  $\delta\underline{B}$  that lead to currents and ion oscillations that damp by Joule and viscous dissipation. Since the generation of  $\delta\underline{B}$  represents work done by the acoustic waves, oblique propagation naturally reduces the damping length generally assumed (cf. III) of the acoustic waves. Nevertheless, we will assume this delineation to be valid in the zeroth approximation. Under these circumstances,  $\underline{B}$  plays an active role only in ED mechanisms and this automatically imposes strong constraints on both the topology of  $\underline{B}$  and its magnitude.

As we will discuss in IV three ED mechanisms (surface waves, reconnection and anomalous Joule heating) require relatively large gradients in  $\underline{B}$  and lead to heating in narrow layers. Since large gradients in  $\underline{B}$  implies large current densities we are faced with not only understanding the ED heating mechanisms but also understanding the global stability of magnetic configurations in which some of these mechanisms are supposed to operate.

Present observational limitations, little can be said of the actual internal structure of the magnetic configurations observed. However, there are few theoretical points that can be made.

The first point is that steady state heating by definition is impossible in force free fields ( $\underline{J} \parallel \underline{B}$ ), because the dissipation of magnetic free energy must be balanced by a divergence of a Poynting flux, which by definition is zero for a force free field. Hence, coronal heating by an ED mechanism which operates in a force free

field, will be time dependent because it must occur in a time at least shorter than the radiation time scale.

The second point is that any ED coronal heating mechanism that dissipates magnetic free energy must heat the plasma in a time  $t_q$  shorter than the characteristic plasma cooling time  $\tau_c$ . This implies that the growth time  $t_g$  of the mechanism must be shorter (generally  $t_g < t_q$ ), which for the anomalous Joule heating and reconnection mechanisms requires strong magnetic fields and/or strong magnetic field gradients (cf. IV). However, strong fields and field gradients are difficult to maintain in the open corona without external stresses due to the tendency of the field to expand so as to relax the gradients and weaken the field. Hence, to heat the corona either in steady state or quasi-steady state by an ED process such as anomalous Joule dissipation it is necessary for the divergence of the Poynting flux to resupply energy in a time shorter than a plasma cooling time (by conduction, radiation or expansion) and also with a magnitude sufficient to compensate both for energy dissipated by the ED mechanism and the energy lost to expansion. Such constraints impose further constraints on the current generators which are necessary for ED mechanisms to function.

In principle the whole current system can be defined by a knowledge of the velocity field, magnetic field and the pressure gradients. This can be seen by deriving  $\tilde{j}_\perp$  from the momentum equation. Neglecting gravity one obtains:

$$\tilde{j}_\perp = \frac{c \tilde{B} \times \nabla p}{B^2} - \frac{c \rho}{B^2} \frac{d\tilde{v}}{dt} \times \tilde{B}. \quad (4)$$

From  $\nabla \cdot \underline{j} = 0$  one finds  $\underline{j}_{||}$  after integrating along a field line:

$$\underline{j}_{||} = -B_0 \int_0^S \left\{ \frac{sc}{B} \frac{d}{dt} \left( \frac{\underline{B} \cdot \nabla \times \underline{v}}{B^2} \right) + \frac{\underline{j}_\perp \cdot \nabla \underline{B}}{B^2} - \frac{1}{sB} \left[ \underline{B} \times \left( \frac{s}{B^2} \frac{d\underline{v}}{dt} \right) \cdot \nabla \rho \right] \right\} ds. \quad (5)$$

Thus a solution to the appropriate equations defining  $\underline{v}$ ,  $\underline{B}$  and  $\rho$  together with the observable boundary conditions on  $\underline{v}$ ,  $\underline{B}$  and  $\rho$  at the photosphere, do specify  $\underline{v}$ ,  $\underline{B}$  and  $\rho$  in the solar atmosphere and thus  $\underline{j}_{||}$  and  $\underline{j}_\perp$ .

### III. HEATING BY ACOUSTIC PROCESSES

The theory that has been popular for more than 20 years is that the chromosphere and the corona are heated by periodic shock waves generated by the turbulent motions in the upper layers of the convection zone. The theory is based on the pioneering work of Biermann (1946), Schwarzschild (1948) and Schatzman (1949) and later extended by De Jager and Kuperus (1961), Osterbrock (1961), Uchida (1963), Bird (1964), Ulmschneider (1967), Kopp (1968) and Kuperus (1965, 1969).

The basic idea behind this theory is that due to turbulent pressure fluctuations a field of small amplitude acoustic waves is generated. Due to the exponential decrease of the density the velocity amplitude of the waves as they travel outward increases rapidly according to

$$v(h) \sim v_0 e^{\frac{\mu g h}{2RT}} \quad (6)$$

As the velocity increases non-linear effects of the propagation causes the sinusoidal wave profile to steepen into a saw tooth like sequence of shocks. In these shocks the dissipation rate is much larger than in the linear acoustic waves. Moreover the heating rate is thus simply determined by the shock wave Mach number,  $M_s$ , and the period  $P$ . In this section we will first discuss the generation of acoustic waves and then the propagation and dissipation of shock waves.

#### 1. Generation of Acoustic Waves

Acoustic waves are generated in the upper convective layers of the Sun by compressional disturbances associated with the turbulent motion. In order to calculate the total amount of acoustic energy leaving the

convection zone solar physicists have profited from the very elegant theory of the aerodynamic generation of acoustic noise by turbulence developed by Lighthill (1952, 1954, 1960, 1962).

If  $\rho_1$  is the density fluctuation it follows from the equation of motion and the continuity equation that

$$\frac{\partial^2 \rho_1}{\partial t^2} - v_s^2 \nabla^2 \rho_1 = - \frac{\partial F_i}{\partial x_i} + \frac{\partial^2 S_{ij}}{\partial^2 x_i \partial^2 x_j}, \quad (7)$$

where  $i, j = 1, 2, 3$  indicate the three coordinates  $x, y, z$ ,  $v_s$  is the sound velocity,  $F_i$  are the forces acting on a fluid parcel and  $S_{ij}$  are the turbulent stresses.  $F_i = \rho_1 g \delta_{ij}$  and  $S_{ij} = \rho v_i v_j$ .

It has been shown by Lighthill that there are two equivalent ways of considering the turbulent field. In the first place one may consider the turbulent density  $\rho_1$  at every point as a consequence of the internally acting forces,  $F_i$  and stresses,  $S_{ij}$ . Secondly one may consider the turbulent region a whole subjected to an externally applied system of fluctuating forces and stresses. In both cases equation 7 describes the generation of sound. The sound producing sources are at the right hand side. According to the way the forces and stresses are differentiated we distinguish between *dipole* radiation (the first term) and quadrupole radiation (the second term). The intensity of the sound waves is given by  $I(x) = \rho_0^{-1} v_s^3 \langle \rho_1^2(x, t) \rangle$ , where the brackets denote an averaging over time. The total power output due to dipole radiation is  $P_d = \beta \epsilon M^3$  and due to quadrupole radiation  $P_q = \alpha \epsilon M^5$ , where  $\epsilon = \rho \langle v^2 \rangle^{3/2} L^{-1}$  is the turbulent dissipation,  $L$  is the characteristic turbulent length (e.g. the mixing length) scale,  $M$  is the Mach number of the turbulent motions,  $M = \langle v^2 \rangle^{1/2} v_s^{-1}$ , and  $\alpha$  and  $\beta$  are constants depending on the

turbulent correlations  $\alpha = O(10)$ ,  $\beta = O(0.1)$ . Hence  $P_q/P_d = \alpha \beta^{-1} M^2 > 1$  for  $M > 0.1$  (Unno and Kato, 1963; Unno, 1964; Kuperus, 1965).

Magnetic fields in the turbulent layers may add to the dipole emission (Osterbrock, 1961) as well as to the quadrupole emission (Kulsrud, 1955). An essential difference between dipole and quadrupole radiation is that dipole radiation is strongly peaked in the direction of the forces while the quadrupole radiation is isotropic. According to Kulsrud (1955) the amount of noise generated in the magnetic regions ( $\beta \approx 1$ ) may be an order of magnitude larger than in the non magnetic regions ( $\beta \gg 1$ ). The total flux of energy ( $F_m$ ) can be found by integrating the acoustic power over the depth.  $F_m \sim 10^7 - 10^8$  ergs/cm<sup>2</sup> sec for quadrupole radiation. The uncertainty is essentially caused by the uncertainty in the turbulent velocities (Kuperus, 1965). Remember that  $P \sim v^3$  !, an increase of 10% in  $v$  results in more than a factor 2 in  $P$ .

It has been shown by Kuperus (1972) that this emission can be considerably enhanced in turbulent convection. The quadrupole emission of a turbulent eddy moving with an average convective velocity  $v_c$  is increased by a factor  $(1 - v_c/v_s)^{-6/5}$  in the direction of  $v_c$  caused by the fact that a point overlying the noise generating layer "sees" a relatively larger acoustic eddy when it moves towards it than when it moves away from it.

Moreover, due to the Doppler effect and the existence of a gravitational cut-off frequency  $\omega_1 = \gamma g/2v_s$  for sound waves the transmitted part of the acoustic spectrum is strongly modulated by  $v_c$ . These two effects are taken into account in the above mentioned amplification factor. As a conclusion we may state that it appears likely that a

sufficiently large flux of acoustic waves is generated in the turbulent layers of the upper convection zone and this flux is strongly varying over the surface of the sun depending on the convective state of motion as well as the magnetic structure (supergranulation and network boundaries).

Moreover it is expected that the part of the mechanical energy flux in high frequency sound waves is strongly time dependent since it can be modulated by low frequency oscillations in the upper convection zone and photosphere.

This spatial and temporal behaviour of the acoustic energy flux must be reflected in the structure and development of the overlying layers.

The existence of a gravitational cut-off frequency results in a transmitted and a reflected part of the turbulently generated acoustic waves.

A second cut-off frequency exists the so-called Brunt-Väisälä frequency  $\omega_2 = (\gamma - 1) g v_s^{-1}$  which is slightly smaller than  $\omega_1$ .

Waves with  $\omega < \omega_2$  can propagate as internal gravity waves with small compressive properties in the convectively stable parts of the atmosphere. Inside the convection zone only pure acoustic waves can be excited (Moore and Spiegel, 1964; Moore, 1967). However, the gravity waves can be excited by overshooting from the convection zone. The spectrum of these waves has been calculated by Stein (1968). Both the acoustic and the gravity modes are refracted in the upward direction because the temperature first decreases. After the temperature minimum is reached, the waves have the tendency to refract away from the radial direction so that a possibility exists for trapped waves in the chromosphere. Uchida (1965, 1967) demonstrated that the transmissivity of gravity waves is several orders of magnitude smaller than for acoustic

waves. Indeed the observed oscillatory character of the photosphere confirms the existence of evanescent wave modes peaked around a period of 300 sec. with strong acoustic properties (Deubner, 1975; Stein and Leihbacher, 1976; Ulrich, 1977; Gough, 1978), while little energy is contained in the gravity modes. Moreover it has been suggested by Souffrin (1966) and Schatzman and Souffrin (1967) that the low frequency gravity modes are strongly damped by radiation. It is thus unlikely that these waves play a significant role in directly heating the corona as has been suggested by Whitaker (1964), although they may be present in the lower chromosphere. It should be mentioned here that due to the rotational character of these waves an interesting possibility exists for coupling with Alfvén waves (Lighthill, 1967). In that case the magnetic field may pick up part of the gravitational wave energy and transport it with much less losses through the chromosphere and the transition region. The propagation of magnetohydrodynamic waves will be discussed in the context of the electrodynamic processes.

## *2. Propagation and Dissipation of Shock Waves*

As the acoustic waves travel upwards their amplitude increases due to the decreasing density. Therefore non linear effects occur deforming the wave profile. The velocity of a point in the wave profile depends on the particle velocity  $v$ .

Because of the difference of the velocities between the crest and the trough of a sinusoidal wave, the distance traversed by a wave of wavelength  $\lambda$  after which the crest and the trough coalesce is  $d = v_s \lambda / 2v(\gamma+1)$ . It can be shown that shock waves are formed after three to four scale heights (Kuperus, 1965; 1969).

During the propagation of a shock wave through the atmosphere some of the shock wave energy is transferred into thermal motions due to

viscous processes. This shock wave dissipation is a direct consequence of the entropy jump through a shock transition. This entropy jump is uniquely determined by the shock wave Mach number  $M_s$  or by the shock strength  $S = \Delta p/p_1$  where  $\Delta p$  is the pressure jump and  $p_1$  the pressure in the preshock medium.

In a quasistationary atmosphere permeated by a sequence of shock waves separated by the accompanying system of expansion waves the energy rise after a shock transition is balanced by non adiabatic processes such as radiation and heat conduction during the expansion.

It then becomes relatively simple to calculate the heating during one cycle as has been outlined by Schatzman (1949). The amount of heating during such a quasi stationary cyclic process

$$Q_R = \oint p d(p^{-1}) \quad (8)$$

is immediately connected to the net increase in enthalpy,  $h$ , after the medium has been expanded to its preshock conditions. The net heating rate is then  $E_h = \rho \Delta h / F$ . For weak shock waves the heating function is given by

$$a_{T_R} \sim \frac{2}{3} \frac{\rho v_s^3}{P} \frac{(M_s^2 - 1)^3}{(\gamma + 1)^2}, \quad (9)$$

In general the enthalpy jump is a complicated function of  $M_s$  given by the Rankine Hugoniot relations. The propagation of a sequence of shocks in a stratified medium has only been treated in an approximate way.

(Kuperus 1965; Ulmschneider 1967, 1970, 1971, 1974; Kopp 1968; Flower and Pineau des Forêts 1977).

If one assumes zero net-displacement, the energy in a shock wave  $D$

can be written as

$$D = [\rho] \cdot [\nu] \cdot t_0 \quad (10)$$

where brackets denote the jump through the shock front and

$$\frac{1}{\tau_0} = \int dt \frac{p - p_1}{p_2 - p_1} \cdot \frac{\nu}{\nu_2} \quad (11)$$

is the so-called time-energy integral (Gonczi et al., 1977). Brinkley and Kirkwood (1947) calculate the shock propagation by assuming  $t_0$  proportional to the characteristic time for the wake

$$\tau^{-1} = \frac{\partial \ln \{ (p - p_1) \nu \}}{\partial t} \quad (12)$$

In this case exact differential equations can be obtained for  $dt_0/dx$  and for  $dD/dx$ . An immediate consequence of the Brinkley Kirkwood theory is that  $dt_0/dx > 0$ , which means that the wake of a shock increases as the shock propagates. It thus appears that if shock waves are generated periodically (e.g. by steepening of a sinusoidal sound wave with period  $P$ ), a subsequent shock catches up the wake of the preceding shock, so that  $t_0 \approx P$  and therefore shocks cannot be considered as separated shocks.

Actually each shock evolves in the wake of the preceding shock. In ... it is customary to assume that  $t_0 \approx \text{constant}$  so that one can derive an implicit relation for the shock Mach number

$$\frac{dM_s}{dx} = G(M_s, \rho_1(x), p_1(x)), \quad (13)$$

where  $\rho_1$  and  $p_1$  are the density and the pressure of the undisturbed atmosphere through which the shock train travels.

Another way to treat shock wave propagation is to replace the atmo-

sphere by a great number of discrete layers and calculate the transmission through each contact discontinuity. This is the Chisnell method (Chisnell, 1955) that is only applicable if the reflected shocks are neglected.

The third method that is fruitfully applied is the so-called Whitham method (Whitham, 1957) which is a modification of the method of characteristics and essentially means the substitution of the shock relations into the equations of motion in characteristic form applied immediately behind the shocks. It is essential for Whitham's method to apply the characteristic relation along the slightly different shock direction. Intuitively it comes close to neglecting the negative characteristics. The accuracy depends on the degree of curvature of the positive characteristics. Extensive discussions of Whitham's method as well as exact expressions for the shock strength gradient are given by Stefanik (1969) and Kopp (1968), who applied this method to the heating of the solar corona. The reader is also referred to Gonczi et al. (1977) who question the validity of this approximation particularly in the case of periodically generated shock waves for which  $t_0 = P$ . Gonczi et al. (1977) give a self consistent formulation of a shock heated atmosphere. For a shock strength  $S$  the maximum flux consistent with  $S$  is reached when the shocks overlap. On the other hand a given flux of shock waves,  $F$ , determines a minimum value for  $S$ . Applying this to the upper chromosphere results in strong shock waves ( $S > 10$ ). Their relation between  $F$  and  $S$  is

$$F \sim \frac{p T^{1/2} S^2}{\left(1 + \frac{S}{2}\right)^{3/2} \left(1 + \frac{\gamma+1}{2\gamma} S\right)^{1/2}}, \quad (14)$$

where  $p$  and  $T$  are the average pressure and temperature. It follows

from (9) that in the lower transition region where the temperature scale height is much less than the dissipation length for the shock waves so that in a first approximation  $\mathcal{F} \approx \text{constant}$  the shock strength rapidly decreases with height until  $S \lesssim 1$ . This is a direct consequence of the reflections of the waves against the temperature gradient. The atmosphere is thus roughly divided into three layers:

1. the chromosphere where the acoustic waves steepen into shock waves;
2. the upper chromosphere and the lower transition region which are permeated by strong shock waves.
3. the upper transition region and lower corona which could be heated by weak shock waves though electrodynamic mechanisms are likely more relevant as will be discussed in IV.

It should be noted that the upper chromosphere and the lower transition region are just the sites of the most violent motions. (See II)

Large amplitudes are confined to a narrow dynamical boundary layer.

The acoustic theory seems to be applicable to the lower parts of the outer atmosphere, e.g. chromosphere and transition region outside the magnetic concentrations.

#### IV. HEATING BY ELECTRODYNAMIC PROCESSES

##### IV.1. Overview

IV.1.1. GENERATION OF ELECTRODYNAMIC ENERGY Electrodynamic processes can play a critically *active* role in the upper chromosphere, transition layer and corona where supposedly  $\beta < 1$  as well as in the magnetic concentrations throughout the whole atmosphere, (J. Hollweg, 1979).

The generation of electrodynamic energy in the solar atmosphere arises when plasma is forced to flow across magnetic field lines. These electric fields stem from, for example,  $\vec{v} \times \vec{B}$  terms in Ohm's law and represent the electric field in a frame that is at rest with respect to the plasma which is moving at velocity  $\vec{v}$ . In addition, the electrons and ions which comprise this moving plasma respond differently when under the influence of magnetic and electric fields. The net result is that local charge separation can occur in the frame of the moving plasma, which, upon neutralizing, results in the generation of electrical currents. These currents then modify the ambient magnetic field. Thus, a simple process results in the conversion of mechanical energy associated with velocity fields into electrodynamic energy associated with the production of electric and magnetic fields. These processes are particularly important in the photosphere and below where a large reservoir of mechanical energy exists (cf. III). Therefore, to understand the generation of electrodynamic energy it is essential that observations focus upon both the power spectrum and polarization of velocity fields *with respect to the magnetic field*.

IV.1.2. PROPAGATION OF ELECTRODYNAMIC ENERGY Under certain conditions to be discussed below, electrodynamic energy can propagate away from

its site of generation represented by an electrodynamic Poynting flux. The destiny of this energy also depends upon a number of conditions, but it is possible to isolate two possible courses of evolution. Specifically, these two possibilities are irreversible dissipation into thermodynamic endproducts and in-situ storage as magnetic energy. In this regard, there are also two different ways that the solar atmosphere can be coupled to the photospheric driver. On the one hand, the atmosphere could be driven in a time-averaged steady state -- the response depending actively upon the details of the photospheric driver (e.g. the power spectrum and polarization of the photospheric velocity fields). On the other hand, however, the response might depend only passively upon the details of the driver. For example, in the case that the photospheric driver could trigger a transient unloading of stored magnetic energy. These points are readily illustrated by Poynting's theorem which is given by:

$$\vec{j} \cdot \vec{E} = -\frac{c}{4\pi} \nabla \cdot (\vec{E} \times \vec{B}) - \frac{1}{8\pi} \frac{\partial B^2}{\partial t}, \quad (15)$$

where the first term on the right hand side represents the net flow of energy into or out of a control volume (i.e., the divergence of the Poynting flux) and where the second term represents a gain or loss of magnetic energy in this volume (See also equation 16 and following discussion). The first term clearly represents a coupling between regions of electrodynamic energy generation such as the photospheric driver and external region such as the corona. As mentioned above, this coupling could very well be in a time-averaged steady state. The second term on the right hand side of equation (15) is explicitly time dependent and thus represents the storage of magnetic energy within the corona.

In order that the corona be driven in a time-averaged steady state, it is essential that the generated electric fields be at least partially in phase with and have a component that is perpendicular to the generated magnetic field. Since these fields stem from  $\underline{v} \times \underline{B}$  terms in Ohm's law, a steady state energy input is always associated with mass flows across magnetic field lines, both at the photospheric driver and within the solar corona. It is important to note that this is a fundamental law of magnetohydrodynamics and is completely model independent.

IV.1.2. DISSIPATION OF ELECTRODYNAMIC ENERGY The third aspect of electrodynamic coupling of the photosphere, chromosphere and corona is the dissipation of electrodynamic energy. Ionson et al (1980) point out that this is conveniently illustrated by the following equation for  $\underline{j} \cdot \underline{E}$ , which is obtained by using Ohm's law, the one-fluid momentum equation and Maxwell's equations:

$$\underline{j} \cdot \underline{E} = \eta_{\perp \text{ electric}} \underline{j}_{\perp}^2 + \eta_{\parallel \text{ electric}} \underline{j}_{\parallel}^2 + \eta_{\text{mechanical}} \underline{j}_{\perp}^2 \quad (16)$$

The first term on the right hand side of equation (16) represents electron "Joule" heating by cross-field currents. The generalized resistance  $\eta_{\text{electric}}$  includes both anomalous as well as classical transport processes. The second term also represents electron "Joule" heating, but in this case it results from the dissipation of field-aligned (i.e. force-free) currents. Here too,  $\eta_{\text{electric}}$  includes both

anomalous and classical transport phenomena. The third term represents viscous heating and mechanical acceleration of the ions due to both anomalous and classical processes (i.e.  $\eta_{\text{mechanical}}$  includes these microscopic effects). In this case viscous ion heating results from sheared ion velocity fields that are driven by Lorentz forces.

#### IV.2. Heating by Magnetohydrodynamic "Body" Waves

Magnetohydrodynamic body waves are well defined *normal* mode oscillations which transport energy via the intrinsic elasticity and tensile properties of the magnetoplasma. These waves are characteristic of a uniform plasma and thus the addition of plasma non-uniformities plays only an extrinsic role (e.g. reflection and refraction phenomena). There is also a second-class of waves which exist only in the presence of non-uniformities. These so-called "surface waves" will be discussed in section IV.3. With regard to hydromagnetic body waves there are three different wave modes, i.e. Alfvén, fast and slow modes, whose dispersion relations are given by:

$$\begin{aligned}\omega_A^2 &= k^2 v_A^2 \cos^2 \theta \\ \omega_F^2 &= \frac{k^2 v_A^2}{2} \left[ (1+\beta) + \left[ (1+\beta)^2 - 4\beta \cos^2 \theta \right]^{1/2} \right] \\ \omega_S^2 &= \frac{k^2 v_A^2}{2} \left[ (1+\beta) - \left[ (1+\beta)^2 - 4\beta \cos^2 \theta \right]^{1/2} \right]\end{aligned}\quad (17)$$

where  $\theta$  is the angle between  $\underline{k}$  and  $\underline{B}$ .

In equation (17)  $\omega_A$  represents the eigenfrequency of shear Alfvén waves which are independent of magnetic and thermal compression and depend solely upon the tensile properties of the magnetoplasma. The eigenfrequencies  $\omega_F$  and  $\omega_S$  represent "fast modes" in which the magnetic and thermal compressions are in phase and "slow modes" in which the magnetic and thermal compressions oppose one another.

Before the advent of ATM instrumentation on SKYLAB and the OSO series of satellites, the solar corona and in particular coronal active regions were observed to be hot but lacking in any internal fine structure.

Thus, the theorist was free to treat an unstructured atmosphere. As

such their investigations focused upon hydromagnetic body waves.

Motivated by the qualitative expectations of Alfvén (1947), Giovanelli (1949), Van de Hulst (1953), Piddington (1956) and De Jager (1959), Osterbrock (1961) provided the first quantitative study of the generation, propagation and dissipation of hydromagnetic body waves in a structureless, plane-parallel solar atmosphere. The essence of Osterbrock's model was that fast-modes are generated in the highly turbulent convection zone (see section III). These waves propagate upwards into

the chromosphere where steepening into shocks occurs just like the acoustic waves. Because of their small scale size, these shock waves are efficiently dissipated within the chromosphere and play an important role in the heating process. Some of the fast modes, however, may bleed into the corona. Upon entering the corona, these fast modes couple to pure Alfvén waves which, because they are non-compressive, are weakly damped and therefore have quite large damping lengths. The problem with this scenario is that it depends quite strongly upon the details of the ambient plane-parallel chromosphere and corona. Specifically, the variation in the plasma  $\beta$  with height above the convection zone critically determines the applicability of Osterbrock's approach which invokes geometric wave optics. Thus the merit in Osterbrock's work is in his discussion of relevant dissipation processes. In particular, Osterbrock focused upon *collisional* dissipation processes (i.e.  $\tau_{\text{wave}} > \tau_{\text{collision}}$ ) for both  $\beta > 1$  and  $\beta < 1$  regions. For example, the damping lengths for Alfvén, fast and slow modes in a  $\beta > 1$  region of the atmosphere are given by:

$$L_{\text{Alfvén}} = \left[ \xi_{\text{Joule and ion viscosity}}^{-1} + \xi_{\text{neutral friction}}^{-1} \right]^{-1}, \quad (18a)$$

with  $\xi_{\text{Joule and ion}}^{-1} = \omega^2 \left( \frac{c^2 \eta}{4\pi} + \frac{\eta_{\text{visc}}}{\rho} \right) v_A^{-3}$  and

$$\xi_{\text{neutral friction}}^{-1} = \omega^2 \chi \tau_m v_A^{-3} (1 + \chi)^{-1}.$$

$$L_{\text{Fast}} = \frac{3 v_s^3 \rho}{4 \omega^2 \eta_{\text{visc}}}, \quad (18b)$$

$$L_{\text{slow}} = \frac{3 \chi^3 \nu_A^3}{4 \omega^2 \eta_{\text{visc}}}, \quad (18c)$$

where  $\chi = n_n/n_e$  is the ionisation degree and  $\eta_{\text{visc}}$  the viscosity as defined in section II.1.,  $\omega$  is the wave angular frequency and  $\tau_n$  is the collision time for neutral particles.

For  $\beta < 1$  they are given by

$$L_A = \left[ \sum_{\text{Joule and ion}}^{-1} + \sum_{\text{neutr. friction}}^{-1} \right]^{-1}, \quad (19a)$$

$$L_F = \frac{4 v_{\text{fast}} \tau_0}{S}, \quad (19b)$$

$$L_S = \frac{4 v_{\text{slow}} \tau_0}{S}, \quad (19c)$$

where  $S$  is the shock strength and  $\tau_0$  the shock wave period.

Clearly these damping lengths depend strongly upon the model atmosphere, particularly the magnetic field and the wave period. Thus, we again stress the importance of determining the power spectrum and polarization of velocity fields with respect to the magnetic field as well as a determination of the plasma  $\beta$  as a function of height in the solar atmosphere, both from the point of view of determining damping lengths and transmission coefficients.

As a foundation upon which one can motivate future inquiries, it is possible to establish a range of wave periods  $\tau_{\text{wave}}$  that are relevant to the heating of the chromosphere and corona. Noting from equations (18) and (19) that the dominant damping mechanism for Alfvén waves is Joule dissipation whereas for fast and slow modes it is dissipation through shock formation, we find that  $\tau_{\text{wave}}$  must satisfy

$\tau_w > 3 \times 10^4 B_{\text{chrom}}^{-3/2}$  sec for Alfvén waves and,  $\tau_w > 10^8 v_A^{-1}$  for fast waves otherwise they will be damped before entering

the corona. In obtaining this constraint we have assumed that the scale size of the chromosphere  $l_{\text{chromosphere}} \sim 3 \times 10^8$  cm with a maximum density  $\rho_{\text{chromosphere}} = 8 \times 10^{-8}$  gm/cm<sup>3</sup> and electrical resistivity  $\eta_{\text{chromosphere}} = 3 \times 10^{-12}$  sec. Although we can only guess as to the value of the chromospheric magnetic field, one can make reasonable estimates for both the quiet and the active sun. For example, if we assume that  $B_{\text{chrom}} = 10$  gauss in a quiet region and  $B_{\text{chrom}} = 300$  gauss in an active region, it follows that only Alfvén waves with  $\tau_w > 16$  min couple to the quiet corona whereas waves with  $\tau_w \geq 5.7$  sec couple to an active region corona.

Despite these constraints, Habbal et al (1979) have investigated the possibility of heating active regions by fast-modes with  $\tau_{\text{wave}} = 2$  sec. Although such waves are probably dissipated within the chromosphere, Habbal et al's approach might apply for longer period waves provided they assume a *collisional* damping process. Habbal et al assume collisionless transit-time damping which is probably a poor assumption since the collision time is small compared to the wave periods of interest (i.e.  $\tau_{\text{collision}} = 1$  sec  $< 5.7$  sec with 5.7 seconds being the smallest wave period that can penetrate into the corona).

There were two major considerations that motivated subsequent work on hydromagnetic body wave heating of the solar corona. First, there is the problem of invoking the WKB approximation under conditions in which it is not applicable and second, there is the problem of very large damping lengths and thus apparently negligible heating of the coronal plasma by those waves that do manage to penetrate the chromospheric filter. By far, the problem that has received the most attention since Osterbrock's

(1961) initial work has been that associated with mechanisms which increase the damping of Alfvén waves, thereby decreasing the coronal damping length. The motivating factor which led to these studies was the discovery that the corona is in fact quite structured. Thus our previous concept of a corona with a single scale size of the order of  $10^{11}$  cm was completely altered and it became necessary to investigate processes that would result in efficient heating of significantly smaller regions such as low lying coronal loops. Despite the structured nature of the solar atmosphere, many authors neglected variations of the plasma parameters in a direction that is perpendicular to the magnetic field. As such, they assumed that Alfvén waves retained their plane-wave nature (i.e. no cross-field variations in the wave properties). This, however, is a reasonable assumption provided that the WKB approximation is valid.

The problem of studying the propagation of hydromagnetic waves without recourse to the eikonal approximation was not addressed until after Belcher and Davis (1971) observed what appeared to be Alfvén waves propagating in the solar wind with periods of the order of hours. Hollweg (1972, 1978, 1980) recognized the important implications of such observations and was the first to develop a model of Alfvén wave propagation in the solar atmosphere without the restriction of the WKB approximation. Although his analysis disregards various damping mechanisms, it does indicate that long period Alfvén waves of the order of hours can propagate through the chromosphere with sufficient power ( $P = 10^4 - 10^5$  ergs  $\text{cm}^{-2}$   $\text{sec}^{-1}$ ) to strongly affect the solar wind and heat at least the quiet solar corona. Subsequently, Wentzel (1974, 1976), Uchida and Kaburaki (1974) and Melrose (1977) claimed that Alfvén waves can be converted into fast-mode waves in the corona from bends in the coronal magnetic field. If the scale length of these bends is of the

order of the wavelength of the Alfvén wave then there results a  $\underline{j} \times \underline{B}$  induced density compression which is characteristic of a fast mode. This density compression rapidly leads to steepening of the wavefront into a shock, and subsequent dissipation in about one wavelength. In principle this seems reasonable, but there is a fundamental problem as to how energy can be coupled into the electron and ion microstructure of the plasma. Specifically, the conversion of Alfvén waves into fast modes occurs only when the wavelength of the Alfvén wave is of the order of the loop size. As such, the resulting shock thickness is also of the order of the scale size of the coronal loop and one must be quite careful when investigating heating of the electron and ion populations.

Wentzel (1974, 1976) also invoked weak non-linearities into his analysis of Alfvén waves. Although Alfvén waves are non-compressive to first order in the wave amplitude, they are to second order compressive. Specifically, the second order pressure fluctuation  $P_2 \approx 4\pi^2 B_1^2 / 8\pi$  ( $B_1$  is the first order magnetic fluctuation) and the second order velocity displacement  $v_2 = 4\pi^2 v_A B_1^2 / B_0^2$  ( $B_0$  = ambient field strength) results in a propagating density wave whose intensity is given by  $P_2 v_2$ . Assuming that this density compression couples energy into heat after travelling about one wavelength  $\lambda = v_A \tau_w$  the damping length associated with the driving Alfvén wave is

$$L_{\text{damping}} = (B_1^2 v_A / 8\pi) \times (\lambda / P_2 v_2) = (4\pi^2)^{-2} (B_1 / B_0) (v_A \tau_w).$$

It should be noted that

Wentzel considered only those cases in which the WKB approach was a good approximation (i.e. wavelength small compared to the magnetic field's radius of curvature).

From Wentzel's analysis it is possible to determine what wave periods are needed in order that his mechanism result in efficient heating of the corona. Noting that  $B_1/B_0 = v_1/v_A$  and also noting that the damping length must be smaller than the characteristic scale size of the corona  $\lambda_{\text{corona}} = 1 \times 10^{11}$  cm, it immediately follows from his expression for the damping length  $L_{\text{damping}}$  that the wave period  $\tau_w$  must satisfy  $\tau_w < 3.9 \times 10^{12} (v_1^{\text{max}}/v_A)^2 / v_A$  sec, where  $v_1^{\text{max}}$  is the maximum observed velocity displacement in the solar corona. Since  $v_1^{\text{max}} \approx 2 \times 10^6$  cm/sec (Berring and Feldman, 1974), we find that only waves with  $\tau_w < 24$  min can efficiently heat quiet regions whereas only waves with  $\tau_w < 5.7$  sec can heat active regions. Although Wentzel's damping mechanism is consistent with the heating of quiet regions there is some question as to its applicability in active regions since only waves with  $\tau_w > 5.7$  sec escape from chromospheric damping. It should also be pointed out that despite the large wavelengths associated with wave periods  $\tau_w < 24$  min (i.e.  $\lambda < 3 \times 10^{10}$  cm) Wentzel's analysis is still approximately valid since quiet regions are inherently much less structured than active regions. It would, however, be more accurate to develop a model of quiet region heating that incorporates the virtues of both Hollweg's (1978) non-eikonal method of studying long period Alfvén waves and Wentzel's (1974, 1976) non linear damping mechanism.

### 3. Heating by Magnetohydrodynamic "Surface" Waves

Ionson (1977, 1978, 1980) was the first to recognize the importance

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of global structure oscillations in the inherently structured active region corona. Ionson, following Chen and Hasegawa (1974) and Hasegawa and Chen (1976), has stressed that the response of a plasma structure to external driving is *qualitatively different* from that which occurs in an unstructured setting. This difference stems from the fact that a plasma structure hosts a continuous spectrum of internal oscillators (i.e. hydromagnetic body waves), each representing the oscillatory characteristics of an infinitesimally small piece of the plasma. For example, consider a coronal loop which is characterized by a parabolic depression in the Alfvén speed across the minor radius of the loop (i.e.  $v_A = v_A(x)$ ). For a given driving period  $t_D$ , there will be a continuous spectrum of excited Alfvén waves whose wavelengths satisfy  $\lambda(x) = v_A(x) t_D$ . The existence of such a continuous spectrum of Alfvén waves (other body waves could be excited as well) results in both a "global" and "local" response of a nonuniform plasma structure to external driving. The global response reflects the integrated response of all the internal oscillators which in this case are Alfvén waves. As such, it is associated with a spatially coherent oscillation of the entire plasma structure. The local response, on the other hand, corresponds to that of a single internal Alfvén wave and is essentially a single field line oscillation. Ionson (1977, 1978) originally called these global structure oscillations Alfvénic *surface* waves to stress that the presence of *nonuniformity* is essential to their existence (c.f. Wentzel 1978, who discusses other surface waves). It is important to note that a discontinuous surface is not required. The wavelength of the excited surface wave is given by

$$\lambda_A = \frac{1}{\tau_w} \left[ \frac{B_{\max}^2 + B_{\min}^2}{4\pi(\rho_{\max} + \rho_{\min})} \right]^{1/2} \quad (20)$$

and in essence represents the field-aligned auto-correlation length associated with the flow of electrodynamic energy. Ionson (1977, 1978) has also pointed out that standing surface waves could result provided  $\lambda_p = n\ell/2$  where  $\ell$  is the loop length and  $n$  is some integer. In addition, Ionson (1980) has pointed out that other global structure oscillations such as magnetosonic cavity modes exist in the natural cavities associated with regions between magnetosonic wave reflection points. In particular, multiply reflected magnetosonic body waves phase mix and establish a discrete spectrum of standing (in the direction of the non-uniformity) magnetosonic cavity modes. In this regard, Ionson (1980) has proposed that coronal structures be interpreted as either *hydromagnetic resonance cavities* such as coronal loops or *hydromagnetic waveguides* such as coronal holes. This viewpoint has also been discussed by Querfeld and Hollweg (1980).

One of the most interesting aspects of global structure oscillations such as Alfvénic surface waves and magnetosonic cavity modes is their *linear* coupling to the Alfvén wave continuum at local Alfvén resonances where for example  $\lambda(x=x_{\text{resonance}}) = \lambda_A$ . This feature was noted by Ionson (1977, 1978) in connection with the linear degradation of Alfvénic surface wave energy into localized body wave energy at a spatial location within a coronal loop. The so-called process of "resonance absorption" (Ionson 1978, and references therein) is a specific, but extremely important example that has been found to control the transfer of energy from an external MHD "driver" to the interior of the "driven" plasma structure. In simple, qualitative terms, this process is associated with the existence of an Alfvénic resonance within the plasma structure where the local impedance (i.e. related to the local Alfvén frequency) matches that of the external driver (i.e. related to

the oscillation frequency of the spatially coherent external driver). At the spatial resonance, where impedance matching occurs, dissipation processes, such as ion viscosity, play a critical role in converting the incoming Poynting flux of electrodynamic energy into heating the coronal plasma. Since the dissipation rate increases with decreasing scale size and since the resonant absorption layer is naturally small compared to the global system, it follows that plasma heating and acceleration is localized within the resonance region. In this regard, the resonance absorption process facilitates the dissipation of large scale mhd disturbances by providing a "sink" (i.e. the resonance absorption layer) for large scale electrodynamic energy which, because of its inherently small spatial scale, results in an efficient conversion of this energy into thermodynamic endproducts.

It is important to note that only within the resonance absorption layer, will the effects of local irreversibilities make an explicit showing. That is to say, they will define the spatial extent of the resonance absorption region and determine the local amplitude of the fluctuating electrodynamic fields.

Thus, the relevance of Ionson's work is in his pointing out the existence of large scale structure oscillations such as the Alfvénic surface waves and his discussion of how the energy associated with these large scale oscillations is coupled into the microstructure of the plasma, i.e., the resonance absorption process. Although Ionson did not self-consistently include all aspects of generation, propagation and dissipation of electrodynamic energy into his analysis, he did underline the potential importance of 300 second oscillations in heating solar coronal loops.

#### IV.4. Heating by Currents

IV.4.1. CORONAL HEATING BY ANOMALOUS JOULE DISSIPATION Over the last few years numerous proposals have been made suggesting that coronal heating may result from the dissipation of currents (Tucker 1973; Nolte et al 1977; Rosner et al 1978; Hinata 1979, 1980). There are basically three physical mechanisms capable of dissipating such currents: Joule dissipation, reconnection and double layers. Joule dissipation and reconnection are closely related, differing only in that pure Joule heating causes no topological changes in the magnetic flux surfaces of the magnetic structure, while reconnection causes such topological changes together with Joule heating and strong convective flows. Here we will review critically the possibility that coronal heating may result from either Joule heating or reconnection since double layers have not, as yet, been proposed as a coronal heating mechanism (for a review of these three mechanisms c.f. Spicer and Brown 1980).

To develop a theoretical framework in which we can study the feasibility of either Joule heating or reconnection as viable coronal heating mechanisms we must specify the physical constraints on such heating mechanisms imposed by both theory and observations. The two most obvious constraints are that in *steady state*: (i) the local power input from the heating mechanism  $\epsilon_H \geq \epsilon_L$  (the local losses) and (ii) the power input from the heating mechanism integrated over the fractional volume of the magnetic structure that is heated must roughly equal the integrated power loss from the entire volume of the magnetic structure.

$$\int_{\Delta V} \epsilon_H dV = \int_V \epsilon_L dV, \quad (21)$$

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where  $\Delta V$  is the fractional volume in which the heating mechanism operates and  $V$  is the total volume of the structure. Obviously if heating is occurring uniformly throughout the volume and  $\epsilon_H \approx \epsilon_L$  at each point then  $\Delta V = V$ . However, both reconnection and Joule heating require that  $\Delta V/V \ll 1$  so that  $\epsilon_H \gg \epsilon_L$ . As we will see these requirements fundamentally alter the energy balance arguments previously proposed (e.g. Rosner et al 1978b).

To demonstrate that both Joule heating and reconnection require  $\Delta V/V \ll 1$  and thus  $\epsilon_H \gg \epsilon_L$  we note that at the very *least* the heating mechanism must overwhelm radiation losses,  $\epsilon_L = n^2 f(T)$  (see section II).

For irreversible heating both Joule heating and reconnection result in  $\epsilon_H = \eta J^2$ , where the resistivity  $\eta = 4\pi v_{ei} \omega_{pe}^{-2}$ ,  $v_{ei}$  being the effective electron-ion collision frequency,  $\omega_{pe}$  the plasma frequency,  $J = -ne \underline{V}_D$  and  $\underline{V}_D$  being the drift velocity of the electrons relative to the ions. Using  $\epsilon_H \gg \epsilon_L$  we find

$$v_D^2 \gtrsim f(T) / e^2 \eta \quad (22)$$

Taking (22) as an equality and substituting both the expressions for  $\epsilon_L$  and  $\epsilon_H$  into (22) we find  $\int_{\Delta V} n^2 f(T) dV \approx \int_V n^2 f(T) dV$ , which is satisfied if  $\Delta V = V$ . However, by using Ampere's equation  $\nabla \times \underline{E} = (4\pi/c) \underline{J}$  we find that the characteristic width of the current channel satisfying (22) is

$$\delta l \lesssim \frac{c \delta B}{4\pi \eta} \sqrt{\frac{\eta}{f(T)}}, \quad (23)$$

where  $\delta B$  is the change in the current produced magnetic field across the current channel. If we take the heating to occur along the whole length  $L$  of a cylinder of radius  $r_0$  then the ratio of  $\Delta V/V \approx \delta l / r_0$ ,

which is usually much less than one. To make  $\Delta V/V \approx 1$  would cause the magnetic field associated with the current density required by (22) to be  $\geq 10^3$  g in the corona. Hence, we arrive at a contradiction so that  $\Delta V/V \ll 1$  and thus  $\epsilon_H \gg \epsilon_L$ , from which we must conclude that  $V_D^2 \gg F(T)/\eta e^2$ . However, there is a limit to how high  $V_D$  can become before exciting various micro instabilities that are capable of increasing  $\eta$  by several orders of magnitude. When  $V_D$  reaches a magnitude roughly comparable to or greater than the ion sound speed  $C_s = (k_B T_e/m_i)^{1/2}$  these instabilities are excited and one has to take into account the phenomenon called "anomalous resistivity" (section II.3). Under these circumstances the Joule heating becomes anomalously large so that  $\epsilon_H \gg \epsilon_L$  is easily satisfied. Anomalous Joule heating is the basis for the coronal heating models proposed by Tucker (1973) and Rosner et al (1978) and examined in greater detail by Hinata (1979, 1980) and Vlahos (1979).

We stress that an ideal force free field is incapable of supporting such an anomalous heating since for more than a dissipation time force free fields have zero Poynting vectors and thus  $(8\pi)^{-1} \partial B^2/\partial t = -\underline{J} \cdot \underline{E}$  (see eq. (1)). Using  $\underline{E} = \eta \underline{J}$  (Ohm's Law) and  $\underline{J} = (\alpha c/4\pi) \underline{B}$  for a force free field, it follows that  $B^2 = B_0^2 \exp(-t/t_0)$  where  $t_0 = 4\pi/\alpha c^2$  and  $\alpha = \nabla \times \underline{B}/B$ . Thus if the current associated with a force-free field is to drive unstable the ion-acoustic instability (or any other instability that is capable of causing anomalous resistivity) it must have an  $\alpha \approx 1/\delta l$  as defined by (23) on some field line so that the current will only last for a short time. Hence, coronal heating due to anomalous Joule dissipation and driven by a force-free field current is a one shot heating process so that steady state heating by force-free field currents is impossible. Thus, for any semblance

of steady state heating to be possible we must have a non-zero Poynting vector. Further, we can rule out a  $\underline{J}_{\parallel}$  associated with an equilibrium  $\underline{J}_{\perp}$ , since  $\underline{J}_{\perp}$  in equilibrium is determined by the pressure gradient which in turn would necessarily be extremely excessive for coronal conditions if it is to support a  $\underline{J}_{\parallel}$  sufficient to drive a current unstable to microscopic instabilities. Hence, if  $\underline{J}_{\parallel}$  is to drive the current unstable to any instability that causes anomalous resistivity, convective flows must exist in the coronal structures that are generating  $\underline{J}_{\parallel}$ .

Let us assume that such flows do in fact exist so that some semblance of steady state anomalous Joule heating can be achieved. Physically this implies that the drift velocity of the current must satisfy  $V_D \approx C_S$  and the instability will be near marginal stability. Marginal stability is simply a statement of the fact that microscopic mechanisms such as the ion acoustic instability grow very rapidly. The ion acoustic instability growth time is  $t_g \approx m_i/m_e \omega_{pe}^{-1}$  and saturates in  $\approx 20 t_g$  so that if one insists that such mechanisms play important roles in coronal heating the instability must be constantly driven by external means towards instability. Thus for a quasi-steady state to exist the threshold conditions must be marginally satisfied (Schep and Boris 1977; Spicer and Brown 1980). A marginally stable system can pulsate about some quasi-steady state, if the dissipation time is shorter than the build up time. The heat pulses have a temporal width given by the dissipation time while the pulses will be temporally separated by the build up time. However, if the build up time is shorter than the heating time the heating will continue until the dissipation time equals the current build up time because the threshold condition  $C_S$  will increase so that  $\delta \mathcal{L}$  will

decrease until the dissipation time equals the build up time. At this point a quasi-steady state will be achieved with the instability rapidly switching on and off. Averaged over the seeing time of the average solar instrument such a coronal heating mechanism would appear in steady state, so that the assumption that the instability is at marginal stability is a good one if such a mechanism is in fact operating to heat coronal structures.

IV.4.2. HEATING BY RECONNECTION Heating of the corona by reconnection has been proposed by Levine (1973), where he argued that if large numbers of  $J_{\perp}$  driven neutral sheets were to exist in the solar atmosphere the Joule heating and particles accelerated by the reconnection process would serve to heat the corona. His arguments were qualitative in nature, with little support either observationally or from detailed theoretical calculation. Since Levine's hypothesis is basically an application of a flare mechanism to coronal heating, and Rosner et al (1978b) have resurrected the Alfvén-Carlqvist flare Model (Alfvén and Carlqvist 1967; Smith and Priest 1972) as a coronal heating model it would be useful to examine briefly those flare models that use  $J_{\perp}$  driven reconnection by tearing modes (Spicer 1976, 1977; Colgate 1978) to see if they would be practical as scaled down coronal heating models.

Reconnection either by neutral sheets (c.f. e.g. Priest 1976) or by tearing instabilities (Furth, Killeen and Rosenbluth 1963; hereafter FKR) represents a mechanism by which energy stored globally is released locally in a concentrated form. Depending on the collisionality of the plasma and the characteristic magnetic field gradients the reconnection

process can be collisional ( $\gamma_{TM} \ll \nu_{ei}$ ), collisionless ( $\gamma_{TM} \gg \nu_{ei}$ ) or semi-collisional ( $\gamma_{TM} \approx \nu_{ei}$ ), where  $\gamma_{TM}$  and  $\nu_{ei}$  are respectively the tearing mode growth rate and the electron-ion collision frequency. In practice if the plasma is collisionless or semi-collisional the difficulties associated with anomalous Joule heating again arise: that is, extremely steep field gradients and a strong external driver maintaining  $J_{\parallel}$  are again required. On the other hand, the constraints on collisional tearing are not as severe. Nevertheless, regardless of what the collisionality of the reconnection process is there are a number of general arguments that can give us some insight into the viability of reconnection by tearing modes as a coronal heating mechanism.

Reconnection in neutral sheets is treated in steady state, while in more complex magnetic configurations such as loops the tearing mode, a dynamic mechanism, is responsible for the reconnection process. The tearing mode gives rise to *inductive* electric fields parallel to the magnetic field, and it is this electric field that is responsible for either Joule heating and weak particle acceleration (collisional tearing) or strong particle acceleration (semi-collisional and collisionless tearing). This parallel electric field,  $\delta E_{\parallel}$ , is generated within a layer whose thickness  $\Delta$  is small compared to the characteristic scale sizes of the magnetic field configuration  $\delta \mathcal{L}$ . It is within this layer that most of the dynamics of the instability takes place. The first point to remember about the tearing instability is that its wavelength parallel to the magnetic field is infinite, i.e.

$\lambda_{\parallel} = \infty$  ( $\underline{k} \cdot \underline{B} = 0$ ), while its wavelength perpendicular to  $\underline{B}$  must satisfy  $\delta \mathcal{L} / \lambda_{\perp} < 1$  (FKR 1963). The fact that  $\lambda_{\parallel} = \infty$  implies that a significant spatial scale of the magnetic configuration parallel to  $\underline{B}$  is involved in the reconnection process, while a spatial scale

perpendicular to  $\underline{B}$  with a scale size  $\lambda_{\perp} \gtrsim \delta l$  will be involved. Hence, the tearing mode will manifest itself as a long filamentary brightening parallel to  $\underline{B}$  with a thickness  $\lesssim \lambda_{\perp}$ . As a general rule the smaller the ratio  $\delta l/\lambda_{\perp}$  the more effective the tearing mode will be in extracting magnetic energy from global storage and releasing it within  $\Delta$  ( $\Delta \ll \delta l$ ). This follows because wavelengths that satisfy  $\delta l/\lambda_{\perp} \lesssim 1$  strongly bend the magnetic field and thereby generate the largest restoring forces. These large restoring forces in turn lead to very low amplitudes for the modes at saturation. Thus little energy will be extracted from the magnetic field. On the other hand those modes that satisfy  $\delta l/\lambda_{\perp} \ll 1$  generate the weakest restoring forces and grow to large amplitudes before saturation, thereby extracting a larger fraction of the available global free energy supply. However, long wavelength modes not only lead to significant Joule heating (or particle acceleration) but also to strong convection and therefore strong disruptions of the magnetic configuration. Our arguments imply that short wavelength modes ( $\delta l/\lambda_{\perp} \lesssim 1$ ) will not significantly heat the plasma while the long wavelength modes ( $\delta l/\lambda_{\perp} \ll 1$ ) can heat the plasma but will also cause large scale disruptions of the plasma-magnetic field configuration.

There are two types of tearing modes treated by FKR (1963), the slow-tearing modes and the fast-tearing modes. As a rule the thinner  $\Delta$  relative to  $\lambda_{\perp}$  the more likely that the mode will be a fast-tearing mode. The non-linear slow-tearing modes have been treated by Rutherford (1973) and were shown to saturate at low amplitudes, thereby extracting little magnetic energy. However, the fast-tearing mode is known to grow exponentially right up until saturation and for this reason was proposed as an excellent candidate for a flare mechanism (Spicer 1976, 1977). More recently these modes were shown analytically

to be capable of heating the plasma (Spicer 1980). However it was found that only after many e-foldings significant heating begins to occur because the perturbed parallel inductive electric field is very weak initially and must grow over many e-folding times before it is comparable to the electric field driving the equilibrium  $\underline{J}_{\parallel}$ . Since the slow-tearing mode saturates at low amplitudes it is highly questionable whether it should be considered as a good candidate for coronal heating, while the fast-tearing mode which does lead to heating causes strong convective disruptions. In cylindrical geometry these disruptions should appear as a kinking of the cylinder (or loop) (resistive kinks). Unless such behavior is occurring in quiescent coronal loops it is highly questionable whether tearing modes can be considered to play a role in the coronal heating mechanism.

We should also add before closing this section that collisionless and semi collisionless tearing modes are not likely to give rise to scaling laws similar to those obtained by Rosner et al (1978a), because these scaling laws result essentially from balancing thermal collisional conduction against radiation losses. However, since semi-collisionless and collisionless tearing modes result in supra-thermal electrons the typical mean free path of these electrons can easily exceed the scale size of the magnetic configuration itself invalidating the use of thermal conduction as a transport mechanism under such circumstances. Hence, the resulting scaling laws are likely to be quite different.

#### IV.5. A "Unified" Approach to Solar Coronal Heating

Ionson et al (1980) have developed an approach to electrodynamic heating of the solar corona that unifies all previously proposed mechanisms and at the same time self-consistently

incorporates electrodynamic energy generation, propagation and dissipation. Their so-called "transmission line approach" is based upon treating solar plasma structures as either "closed" or "open" transmission lines (e.g. coronal loops and coronal holes) characterized by impedances, Q-values, L/R, RC and LC times. This approach is valuable in that it stresses the fundamental aspects of electrodynamics, viz., inductance, capacitance and "resistance" while at the same time providing a formalism which describes the self-consistent electrodynamic coupling of the photosphere, chromosphere and corona. This coupling is described by a simple circuit equation which includes for currents flowing along and across the magnetic field, i.e.  $j_{\parallel}$  and  $j_{\perp}$  an inductance  $L$ , a generalized "resistance"  $R$  and the plasma capacitance  $C$ . Ionson et al (1980) point out that in certain limits, these equations reduce to those associated with more traditional heating mechanisms. For example, the "quality" or Q-value of the circuit is defined as  $Q = (L/CR^2)^{1/2}$  and is a measure of how good an oscillator the circuit is. If a particular coronal structure is characterized by  $Q \gg 1$  then this structure is a good oscillator and the circuit equation reduces to that associated with hydromagnetic wave-heating theories. In this case, power absorption is highly peaked about the natural oscillation frequency of, for example, a coronal loop. For the most part, the natural frequency of oscillation for most coronal loops appears to be centered around 300 seconds. Therefore, Ionson et al (1980) claim that high-Q coronal loops should be treated as hydromagnetic waveguides. If, however, a particular coronal structure is characterized by  $Q \ll 1$ , then the circuit equation reduces to that associated with "current-heating" theories. In this case, power absorption is *not* highly peaked about some resonance

frequency and indeed, resonances do not even exist.

Ionson et al (1980) also point out that the above theory is subject to potentially *observable* boundary conditions ( $\underline{v}$  and  $\underline{B}$ ) . As such, this approach could be very important in allowing an efficient assimilation of data from ongoing and future solar missions such as the Solar Maximum Mission and Spacelab missions.

## V. CONCLUSION

The mechanism of the heating of the solar chromosphere and corona remains an open question especially if it comes to explain detailed atmospheric structures. The uncertainty is essentially due to our poor knowledge of the magnetic field and the velocity field spectrum with respect to the magnetic field. The magnetic field in the chromosphere and corona appears to be known only qualitatively, while for a heating theory a rather detailed quantitative knowledge of the magnetic field is required. We have emphasized that it is mainly the value of  $\beta$  whose distribution throughout the atmosphere should be known. According to whether  $\beta \gg 1$  or  $\beta \ll 1$  we have divided the existing theories in

- a) the acoustical theories, such as the well-known shock wave theory which has been widely applied also to explain other stellar coronae;
- b) the electrodynamic theories, which may be divided into heating by body MHD waves, surface MHD waves and current heating.

As a conclusion it can be stated that the acoustic theory is losing its range of applicability concerning the expected distribution of  $\beta$ . Figure 1 schematically indicates the range of validity of the acoustic and the electrodynamic theories.

The acoustic theory might be correct with little modifications for most of the solar chromosphere but as soon as the atmosphere shows a high degree of structure such as in the corona and transition layer the magnetic field must play a dominant role.

As far as the electrodynamic theories are concerned it appears that also the current heating theories have a small range of applicability. They are particularly important for flare heating and for flare-like

coronal phenomena.

Most promising are the wave theories. Body as well as surface waves appear to play an important role. In the "open" magnetic structures body waves may be trapped and these structures may act as a kind of wave guide for MHD waves. The "closed" magnetic structures may act as resonant cavities for Alfvénic surface waves and may therefore act as a special kind of absorber for Alfvén waves. We have stressed the importance for a unifying and consistent theory of the formation and heating of coronal structures that eventually bridges the gap between the photospheric "driver" (e.g. the velocity and magnetic fields as they should be observed in great detail in the photosphere) and the observed chromospheric and coronal radiation structures. A first attempt in this direction in which coronal structures are considered as electrodynamic circuits seems to lead to a promising and original way of describing the solar outer atmosphere.

Finally the reader is warned not to reject the acoustic theory categorically, because it is "en vogue" to do so presently, but to consider the ultimate electrodynamic theory as a more powerful theory that eventually includes the acoustic aspects of chromospheric and coronal heating. The prime question which remains unanswered is namely whether the corona is heated by essentially longitudinal waves (acoustic waves and magneto-acoustic waves (fast and slow modes) or by essentially transversal waves such as the Alfvén waves (body and surface waves).

Due to the paucity of the observations of  $\underline{v}$  and  $\underline{B}$  in the solar atmosphere this question can not be settled. The only facts that seem to be generally accepted is that the lower atmospheric oscillations are acoustic in nature and the magnetic fields are strong and con-

centrated in flux tubes in the photosphere. Until we have observed the state of motion of the chromosphere and corona with respect to the magnetic field we have to live with two types of theory, the acoustic theory and the more general electrodynamic theory.

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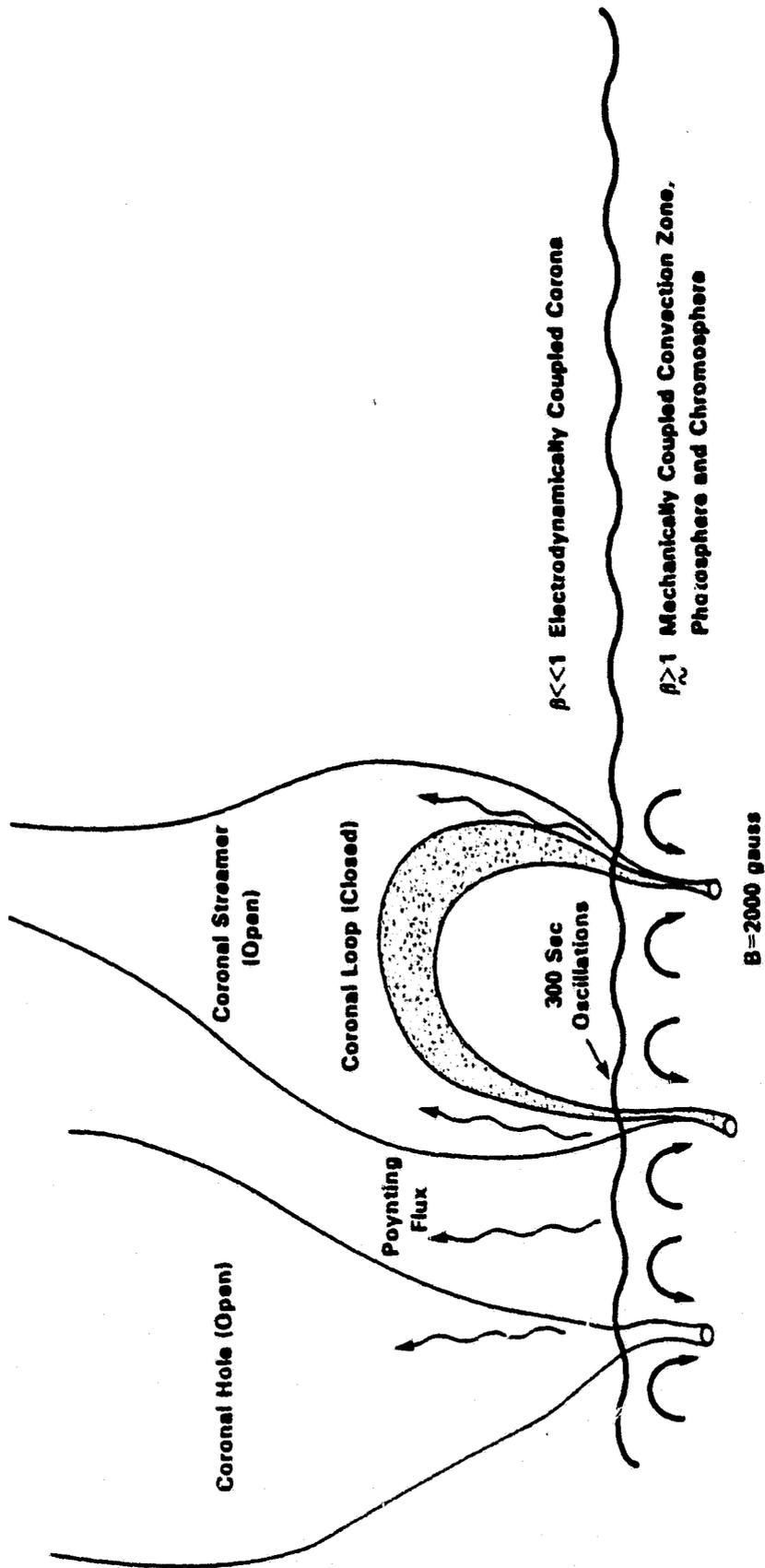
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Table 1.

	Quiet Sun	Coronal hole	Active Region
Coronal energy losses $\text{erg cm}^{-2} \text{sec}^{-1}$	$3 \times 10^5$	$6 \times 10^5$	$10^7$
Chromospheric energy losses $\text{erg cm}^{-2} \text{sec}^{-1}$	$4 \times 10^6$	$4 \times 10^6$	$2 \times 10^7$



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Caption to the figure.

Schematic representation of the magnetic structure of the solar atmosphere. The scale is highly arbitrary. The lower atmosphere where to a large extent  $\beta \gtrsim 1$  is mechanically coupled to the convective and turbulent motions and the 300 sec. oscillations. The corona where  $\beta \ll 1$  is electrostatically coupled to the lower atmospheric velocity field. In the subcoronal layers it is the mechanical energy flux and in the corona the Poynting flux that is responsible for the energy transfer.