ELECTRO MAGNETIC HEATING OF CORONAE

by
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Abstract

Current ideas concerning coronal heating by alternating or direct electric current dissipation are briefly reviewed. It is stressed that Alfven wave heating is not excluded by observations. The promizing aspect of Joule heating by DC current in loops are emphasized, as well as the difficulties still met by such theories.

I) Summary of facts and numbers relevant to the problem

A successful theory of coronal heating should explain : energy, mass and momentum balance of the solar corona as a whole, as well as of specific structures (loops, holes, etc...).

a) Fluxes

Energy budget involves at least the most conspicuous losses, which are radiation, and solar wind enthalpy flux. Figures for these are compiled in the following table.

Region	Loss	Amount (ergs $s^{-1}cm^{-2}$)	Reference
Photosphere	Opt.Rad	$6.41\ 10^{10}$	Bray Loughhead (Table 7.2) Allen
Chromosphere	Opt.Rad UV lines	2-6 5•3 10^6 6•2	Bray Loughhead Vernazza et al (1980) Ulmschneider (1979)
Transition	UV	$6\ 10^5$	Athay-White (1979)
Corona	Opt+UV	$1.5\ 10^4$-$3\ 10^5$	Osterbrock & Athay (1966)
All	X(40-80Å)	$4.8\ 10^3$	Mewe (1978)
"Quiet"	X(2-60Å)	$8\ 10^4$	Vaiana Rosner (1978)
"Active"	X(")	$9.7\ 10^4 \sim 10^5$	Vaiana Rosner (1976)
Holes	X(2-60) Solar Wind	$8\ 10^3$ $7\ 10^5 \sim 10^6$	Vaiana Rosner (1978) Zirker (1980)

Note that coronal holes do not need smaller energy input.

The corona however is strongly structured, and consists of loops of various sizes, density, and physical properties. These loops and structures have widely different energy needs (Vaiana et al. 1976). A comparison of the energy fluxes necessary for each structure stresses again the fact that some of these fluxes differ by a factor as large as 50 from the mean. In the case of active regions it is as large as the flux needed to sustain chromospheric losses.

b) Detailed loop structure and consequences.

Detailed geometrical features of loops have been analysed by Foukal (1975, 1976, 1978), and have shown the following very important properties of cool core coronal loops.

(1) Aspect ratio (radius/length) $\simeq 10^{-1}$, (2) a lower pressure cooler core is imbedded in a hot mantle, (3) most loops last much longer than free fall time and cooling time (they are stationnary), (4) the cool material in the loop center does not suffer a siphon flow from one leg to the other but falls downwards along both legs at a velocity approaching free fall. This special fact implyes that the loop's magnetic field is permeable at least at the top for the matter. In other words matter is not frozen in. It slips through the field. This fact implies the existence of a dissipative process at some portion of the loop length at least.

II) Is mechanical heating by Alfven waves excluded from observation?

Athay and White (1979) have claimed to exclude the mechanical heating theory by putting an upper limit on energy flux from turbulent velocity observed in various chromospheric and transition region EUV line. They state :

$$F_{Mech} \leqslant \rho \xi^2 c_s = \bar{F} \tag{1}$$

and exclude the possibility of higher limits due to Alfven transport on the apparently sound basis that, as you consider an average flux, you should use the average Alfven velocity, i.e. a 1 or 2 Gauss field, and this does not give a higher limit because $\langle v_A \rangle \simeq c_s$. Their upper limit derived from observations is compared to the losses upwards from the corresponding level. It is found to be too small by a factor ten at 2000km.

We believe that their conclusion may be disputed on diverse grounds but we want to stress what we believe is a fundamental difficulty in such estimations. The point is that the mechanical flux may escape observation if it is transported at Alfven speed in concentrated channels, because the more concentrated the field is, the faster corresponding Alfven velocity. We believe that Athay and White made a mistake in evaluating :

$$\langle F \rangle = \frac{1}{S} \iint_S \rho \, dx \, dy \, \xi^2(x,y) \, v_A(x,y) \tag{2}$$

They neglected correlations between ξ^2 and v_A . They state implicitely

$$\langle F \rangle = \rho \langle \xi^2 \rangle \langle v_A \rangle \tag{3}$$

whereas the correct evaluation is

$$\langle F \rangle = \rho \langle \xi^2 v_A \rangle \tag{4}$$

The question is then, if we figure out a composite chromosphere permeated by more intense flux tubes, how much larger should ξ be in the tubes ? A bit of algebra gives the answer $\xi_{Tube} \sim 2.57 \, \xi_{Average}$, but the clumpiness can be arbitrary!!

III) Clues from stellar observations

Giacconi (1980) reports Einstein observations of stellar coronae, and Mewe (1979) reports those of HEAO1. Results of these observation seem to strongly conflict with mechanical heating theories. Recent developments on the observational side, seem to indicate that rotation may be the major parameter which after all determines coronal emission. These aspects are reviewed in RM Bonnet's book (1980).

IV) Alfven waves as a source of heating. Structuring neglected

According to Osterbrock (1961) Alfven waves do not contribute to coronal emission because :

1) They are much less effectively emitted than fast waves.
2) They damp in photosphere-low chromosphere.
3) They are reflected at the transition zone.

All these objections are known to fail. They were due to the fact that Osterbrock pictured out a weak diffuse magnetic field. Let us consider these in turn.

a) Wave emission

Sound wave emission, by a turbulent fluid is calculated by finding the non linear source term to pressure fluctuations which result from otherwise mainly incompressible motions. Because of the isotropy assumption, this source is quadripolar in nature and the result is :

$$\text{Power emitted} = Vol/4\pi \; \rho_0 v^2/(\ell/v) \; v^5/c_s^5 \qquad (5)$$

If gravity is taken into account (Unno, 1964), we also have dipole and monopole sources. Ulmschneider and Bohn (1980) report similar results for a non negligible field. Their result can be expressed as follows for n-pole emission (monopole, n=0)

$$(\text{Power emitted})_{\text{n-pole}} = \text{Vol}\,\frac{\text{En.density}}{\text{Correl.Time}}\left(\frac{\text{Corr.length}}{\text{Wavelength}}\right)^{2n+1} \text{ergs s}^{-1} \qquad (6)$$

The wavelength is that which corresponds to the correlation time. It can easily be checked that Alfven wave emission is favoured in low β conditions.

b) Wave Damping

Linear wave damping proceeds on a scale length :

$$\lambda_{\text{Damp ion}} = v_A^2/\omega^2\,(\eta+\nu)^{-1} \qquad \lambda_{\text{Damp neutral}} = v_A/\omega^2\,\frac{1+\alpha}{\tau_{coll\,i-n}} \qquad (7)$$

Uchida and Kaburaki (1974) have shown the damping to be very weak if B>10 Gauss.

c) Wave reflection and refraction

If we use geometrical optics as a guide, one deduces that slow and Alfven modes are channelled by B ($c_A \gg c_S$) but that pure sound or fast modes, obey Snell's law :

$$\sin i/c_S = C^t \qquad \text{or} \qquad \sin i/c_A = C^t \qquad (8)$$

They refract towards lower velocity regions, and such small fractions as 3 10^{-3} of an isotropic incident flux may reach the corona. The importance of this effect has been stressed by Schatzman (1970), who also discussed the possibility that this dramatic filtering be reduced by non plane layered inhomogeneities and coupling to field channelled modes.

Wave reflection has been discussed by Kaburaki & Uchida, and more recently by Wentzel, Hollweg, B. Leroy (cf. Leroy this meeting).

A recent work by E. Zweibel answers the objection that a fast isotropic MHD mode cannot produce loop type structures. He shows that under the combined effect of thermal instability and refraction, a faint mode heated structure spontaneously develops field aligned fibers.

d) Wave-wave interaction

A full theory of MHD wave heating should describe all the cascades that a wave spectrum suffers being transferred. As a result of these the rate of heating should be predictible. Such strong turbulence theory of heating has not yet appeared for the corona. Dobrowolny et al (1980) however give interesting results for the solar wind turbulence.

The idea of mode coupling has been examined in the weak turbulence approximation (Uchida, Kaburaki, 1974) counterstreaming Alfven waves produce sound waves which dissipate. The process would not work in coronal holes, and is weakly efficient in strong fields (Wentzel, 1976).

V) Role of magnetic structuring

Magnetic structures perturb wave propagation and introduce new modes. The simplest problem is one in which the field structure has one dimension (magnetic wall).

a) Surface waves

It can easily be shown that this structure supports surface waves, i.e. oscillations localized near the interface. No surface wave is absolutely incompressible, but in the limit $k_\parallel \gg k_\perp$ in the surface, this approximation is quite good, and we have a phase velocity intermediate between Alfven waves :

$$\omega^2/k^2 = c_{A1}^2\,\rho_1/(\rho_1+\rho_2) + c_{A2}^2\,\rho_2/(\rho_1+\rho_2) \qquad (9)$$

More complicated dispersion relations, valid in tubes, or when the uncompressibility fails are given by many authors (see for example Roberts and Webb).

None of them has up to now described the generation of surface waves in as much detail as sound emission has been described.

b) Perfect MHD damping of surface waves

The boundary between two regions is not discontinuous. The problem is then to find vibrations in and around a finite current sheet. The wave equation for this problem can be written, and the modes found (Barston, 1964). It is found that they form a continuum between $\Omega(\mp\infty)$ where $\Omega(x)$ is defined by :

$$\Omega(x) = k_{\parallel} v_A(x) \tag{10}$$

Each mode of frequency ω is singular at its resonant point x_ω where :

$$\Omega(x_\omega) = \omega \tag{11}$$

The surface proper mode has disappeared apparently, but it surfaces again if we produce regular packets of Barston's singular modes. It is found that these wave packets, may well start as a typical "surface mode", but that this mode is damped, and leaves assymptotically a secularly vanishing oscillation.

$$\xi(x,t) = t^{-1} \exp(i\Omega(x)t) \tag{12}$$

This decay is easily understood ; it represents a vibration of the structure in which each shell $x = C^t$ oscillates at its own Alfven frequency $\Omega(x)$.
There is then a decay of the oscillations by phase mixing. It has been shown by solving the initial value problem that in the case of finite thickness current sheets, the surface mode appears as an exponentically damped transitory contribution. Sedlacek (1971).

In reality at a high degree of phase mixing, the dissipation by viscosity for example, will be effective, and the damping will be physical, and not only virtual. We may speak of this process as dissipation-less damping, just as we speak of collisionless damping. Dissipation-less damping has been observed in laboratory θ-pinch (Grossman et al, 1972). At a high rate of shearing, the scale of velocity perturbations becomes comparable to the ions gyroradius. Finite Larmor corrections need be considered. In a homogeneous plasma finite Larmor radius corrections to short perpendicular wavelength waves provide the possibility for the Alfven wave to propagate sideways because the dispersion relation becomes :

$$\omega^2 = k_{\parallel}^2 v_A^2 \left(1 + k_\perp^2 \rho_i^2 \left(\frac{3}{4} + \frac{T_i}{T_e}\right)\right) \tag{13}$$

In the inhomogeneous structure problem, this give us finite Barston modes (modified) because the resonantly excited oscillations do not stay where they are excited but propagate sideways. Hasegawa & Chen (1976) have discussed this.

c) Ionson (1978) has pushed this idea up to the construction of an explicit model of loop heating. The most salient feature of dissipation-less damping is that the energy is fed to a small region at the boundary, near the resonance layer. The width of that region is even smaller than the radial wavelength given by Hasegawa's theory if $\omega_{surf\ wave} < \nu_{ii}$. The problem of these sort of theories is to explain how the heat is spread inside the loop. As I understand it Ionson figures out the system as similar to a hot vertical plate in air. A convective boundary layer would develop at its contact, and convect an enthalpy and mass flux upwards. In the loop case there is an extra problem to feature out what the mass and energy become at the top of the loop. Ionson argues that a Rayleigh Taylor instability there diffuses matter in the loop core where it cools by radiation and falls, closing the mass circuit.

VI) Joule heating of the corona

a) Joule heating is a viable theory from global budget point of view

The energy to be released in Joule heating or more generally by interrupting or altering coronal currents is the magnetic free energy :

$$W_M = \int B_1^2/2\mu_0 \; d^3r \tag{14}$$

B_1, the field due to local currents is mainly perpendicular to the current carrying loop, and smaller than a certain fraction of the longitudinal field, otherwise the whole loop turns kink instable ; we find ;

$$W_M \sim 2.5\ 10^{28}\ B_3^2 R_9^4 l_{10}^{-1} \ \text{ergs} \tag{15}$$

The build up time for this energy is :

$$\tau_{feed} \sim 2\pi R_{phot}/v_{Phot} = 6300\ R_{ph8}\ v_{km/s}^{-1} \quad s \tag{16}$$

Hence :

$$\dot{W}_M = 4\ 10^{24}\ v_{km\,s^{-1}}\ R_{Ph8}^{-1}\ B_3^2 R_9^4 l_{10}^{-1} \quad \text{ergs}\ s^{-1} \tag{17}$$

The photospheric field may not have a pressure larger than photospheric pressure, hence $B_3 \leqslant 1.6$. Taking this as a representative value, we find that $\dot{W}_M$ so calculated is large to meet the requirements of X ray luminosity of all known coronal structures. Nevertheless, the idea is not without difficulty.

b) Joule heating is impossible by normal resistivity processes

To convince onself of that, calculate $\langle j \rangle$ from $B\theta$ and use σ from Spitzer's theory and 10^6°K. You find that the current has to be clumped to ridiculous sizes, to increase J^2/σ at the right value. Micro turbulence would occur much before.

c) Dissipation under microturbulent conditions does not work

If microturbulence is to be ever triggered, currents have to be collapsed to a fairly small sheet. Actually, the current density threshold is for ion acoustic or ion cyclotron waves:

$$J^* = 10\, n_e v_{Ti} \sim 16\ n_9\ T_6^{1/2} \quad \text{Amp}\ m^{-2} \tag{18}$$

The total current ($2B_2/\mu_0 l$) at this density has a clumpiness factor very small. It would occupy a strip of width :

$$\delta^* = 500\ R_8\ l_{10}\ n_9^{-1}\ T_6^{-1/2}\ B_3^{-1} \quad cm \tag{19}$$

Once this regions is reached, the evolution is one of marginal stability (Heyvaerts, Kuperus, 1978).

The crucial problem is just to explain the current concentration, why the energy should be fed in the microturbulent sheet, and how the current concentration is maintained.

On all these three crucial points the microturbulence theory of Rosner et al (1978) fails badly in our opinion because :

1) The dissipation interests only a small part of the current structure.

2) The build up of such localized current is incredible. It would imply velocity gradients of 1km/s on 5m in the photosphere.

3) Such small current sheets spontaneously broaden their current profile by tearing modes on very small time scales.

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