CLM-P191

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PLASMA TURBULENCE IN SOLAR FLARES AS AN EXPLANATION OF SOME OBSERVED PHENOMENA

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M. FRIEDMAN
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ABSTRACT

It is suggested that, in Petschek's model of magnetic field annihilation plasma which flows through the boundary layer where its magnetic energy is released is rendered highly turbulent by current driven electrostatic instability. This leads to a physical insight into the mechanism of dissipation, and, by analogy with laboratory experiments on turbulent plasma, can explain the observed X-ray and microwave emissions.

When the microstructure is calculated using electrical conductivity appropriate to highly turbulent plasma, a field configuration exists in which protons can be accelerated to very high energies. The results of some numerical calculations of this process are presented.

It is also suggested that the Type V continuum radio emission may be due to synchrotron radiation from energetic protons, rather than from electrons.

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<u>C O N T E N T S</u>

1. INTRODUCTION

The problem of the rapid annihilation of anti-parallel magnetic fields and the release of their energy in such phenomena as, for instance, a solar flare, seems to be best explained in terms of the propagation of a large amplitude wave or shock, as first proposed by Petschek (1963). It is assumed that large volumes of plasma carrying anti-parallel magnetic fields move towards each other, and that at some points in the common interface local reconnections of the anti-parallel lines first occur by simple resistive diffusion in the plasma, forming magnetically neutral points or lines. The subsequent development of the system is uncertain, although it probably proceeds via some form of resistive instability of the type first discussed by Furth et al. (1963). Petschek's important contribution was to suggest the existence of a steady non-linear solution to the hydromagnetic flow problem in the form of stationary slow shock waves extending from the original "diffusion" region, the initial (upstream) magnetic energy being largely destroyed and converted into particle kinetic energy downstream from the shock fronts, the heated plasma flowing out away from the neutral point region and leaving the flare at the local Alfvén speed V_A. Petschek was able to show that the flow conditions were consistent with the concept of this "switch-off" shock provided the initial flow velocity did not exceed about $0{\cdot}1\ V_A$. An important feature of the model is its remarkable insensitivity to the conductivity of the plasma. The overall time scale for the magnetic energy to dissappear is simply $L/(0{\cdot}1\ V_A) \approx 10^3$ sec if we take for a typical solar flare a characteristic length $L \approx 10^{10}$ cm, $V_A \approx 10^8$ cm sec^{-1}.

Since the original publication of the model there has been little attempt to consider the details of the actual dissipation mechanism, which we feel is important if we are to understand the observed phenomena. This is partly because the physics of the slow shock, which can be shown to require viscous as well as resistive dissipation to be supported, is only imperfectly understood. Although in recent years there has been an appreciable effort devoted to the experimental study of fast hydromagnetic shocks (e.g. Paul et al., 1967), we are unaware of any successful attempts related to the slow shock.

In an earlier note (Friedman and Hamberger, 1968) the authors have shown that the region near the neutral point, whose dimensions are governed only by diffusion processes, would be subject to current induced micro-instabilities (with wave-lengths ≈ Debye distance) and that this would lead locally to a much faster rate of magnetic field diffusion and hence to physically more meaningful dimensions of this region than can be derived on the basis of binary particle encounters alone. This follows from the fact that the effective electrical conductivity in an unstable region can be many orders of magnitude smaller than that usually assumed. As was first suggested by Buneman (1958, 1959) and since confirmed by a number of laboratory experiments (e.g. Fanchenko et al, 1964; Hamberger et al; 1967) the passage of a sufficiently dense current through a plasma leads to turbulence through the extremely rapid growth of plasma oscillations: this serves to obstruct the electron drift motion, and causes the plasma to exhibit an anomalously small conductivity. These experiments usually demonstrate several closely correlated phenomena when the conditions for instability are satisfied, e.g. intense microwave bursts, X-ray emission from the rapidly heated plasma electrons, fast ion production, etc., which have a very strong similarity to the microwave, X-ray, and fast particle emission from solar flares.

It is therefore tempting to explore the possibility that the flare phenomena originate from plasma in a similar state of turbulence to that produced in the laboratory experiments. Since the diffusion region itself is almost certainly too small to account for more than a fraction of the observed radiation, and since we know (from the shock equations) that most of the plasma heating occurs within the boundary layers, we look to a layer itself as a possible turbulent region. Phenomenologically this is reasonable in the light of the experimentally observed strong correlation between the X-radiation (from heated electrons) and the microwave emission (from non-thermal plasma fluctuations). Further, since there is a large magnetic discontinuity across a boundary layer, large currents must flow in this layer and if these are sufficiently dense they will produce plasma instability and turbulence.

We also point out here the existence in the neutral line region of a simple acceleration mechanism which is capable of accelerating a small fraction of the protons entering it to very high energies before they are ejected, and we suggest a physical link between this proton flux and some of the observed continuum radio emission.

2. THE SHOCK WAVE

In the wave model (Petschek, 1963) there is a steady state two-dimensional configuration (Fig.1), in which oppositely magnetized plasma (I) flows into a central region (II) on either side of an x-type neutral point. Diffusion is important to the model only in the region IV close to the neutral point, and the magnetic energy is converted into plasma energy at stationary slow shock fronts III (Kantrowitz and Petschek, 1966). Between these stationary fronts (region II) plasma flows outwards in a direction roughly perpendicular to the initial flow, leaving the flare at the Alfvén speed V_A. The magnetic field changes magnitude and direction at the shock fronts, as shown in Fig.1. In region I the plasma flows with a velocity $u_x = M_o V_A$, where $M_o \lesssim 0{\cdot}1$.

We do not attempt a detailed solution of the structure of the slow shock, for reasons already mentioned, but, using dimensional arguments based on experimental observations, we suggest a self-consistent picture of the shock in which the dissipation is via collective rather than binary interactions.

Regardless of the actual mechanisms we can calculate the change in conditions across the shock from the Rankine-Hugoniot equations. Column (ii) of Table I shows the result of such calculations, the assumed initial conditions of the incoming plasma being shown in column (i). The coordinate system is shown in Fig.1.

TABLE I

SUMMARY OF PLASMA PARAMETERS

Parameter	(i) Upstream	(ii) Downstream	(iii) Principal mechanism responsible	(iv) Characteristic length for discontinuity
B_y(G)	500	0	ohmic (electron scattering)	magnetic diffusion length δ_1
B_x(G)	100	50	-	-
$n(cm^{-3})$	$2 \cdot 10^{11}$	$5 \cdot 10^{11}$		
$-u_x/V_A$	0·1	0·04	viscous (ion scattering)	ion mean free path or ion gyro-radius δ_2
u_y/V_A	0	1		
T(K)	10^4	10^8		$\delta_3 = \max(\delta_1, \delta_2)$

Column (iii) shows the most obvious physical process to account for the change in each parameter, while in column (iv) we suggest a reasonable scale length for each "discontinuous" property. It is not necessary that the scale lengths for the flow and magnetic fields are the same, but it appears more likely that the flow changes within the magnetic discontinuity, i.e. $\delta_2 < \delta_1$ (as in fast shocks with $M > 2 \cdot 5$), than the reverse.

TABLE II

CHARACTERISTIC LENGTHS

	δ (cm)	$j = \frac{c}{4\pi}\frac{B}{\delta}$ (esu)	$v_d = \frac{j}{ne}$ (cm sec^{-1})
λ_{ii}	$10^8 - 10^9$	$10^4 - 10^3$	$10^2 - 10$
ρ_L	$10 - 10^2$	$10^{11} - 10^{10}$	$10^9 - 10^8$
$\frac{c}{\omega_{pi}}$	80	2×10^{10}	$2 \cdot 10^8$
$\frac{c}{\omega_{pe}}$	2	8×10^{11}	8×10^9
$\frac{u}{\omega_{pi}}$	0·2	8×10^{12}	> c ?
$\frac{c^2}{4\pi\sigma u_x}$ (i) classical conductivity	$< 10^{-1}$	$> 10^{13}$	$\gg$ c ?
$\frac{c^2}{4\pi\sigma u_x}$ (ii) turbulent conductivity	30	5×10^{10}	5×10^8

N.B. $\omega_{pi} = \left(\frac{m_e}{m_p}\right)^{\frac{1}{2}} \omega_{pe}$.

Table II lists the possible scale lengths involved, and their associated current densities and drift velocities, which are related by

$$\underset{\sim}{\nabla} \times \underset{\sim}{B} \approx \frac{B_y}{\delta} = \frac{4\pi}{c} \underset{\sim}{j} = \frac{4\pi}{c} ne \, \underset{\sim}{v}_d \quad \ldots (1)$$

Small scale turbulence will be generated by current-driven electrostatic instabilities whenever the electron-ion relative drift velocity v_d exceeds the phase velocity of longitudinal plasma waves (Stringer, 1967; Kadomtsev, 1965) i.e.

$$v_d \gtrsim (kT_e/m_e)^{1/2} \quad \ldots (2a)$$

for Langmuir waves, or

$$v_d \gtrsim (kT_e/m_p)^{1/2} \quad \ldots (2b)$$

for ion sound waves, where kT_e is the mean electron energy, and m_e, m_p are the are the electron and ion masses respectively. When such turbulence exists the effective electron collision frequency is experimentally found (for hydrogenic plasma) (Hamberger and Friedman, 1968) to be

$$\nu_{eff} \approx 10^{-1} \, \omega_{pe} \quad \ldots (3a)$$

if condition (2a) is satisfied, and

$$\nu_{eff} \approx 10^{-2} \, \omega_{pe} \quad \ldots (3b)$$

if the less stringent condition (2b) is fulfilled, where

$$\omega_{pe} = \left(\frac{4\pi ne^2}{m_e}\right)^{1/2} \quad \ldots (4)$$

is the electron plasma frequency for plasma of charge number density n. The corresponding electrical conductivity is related to ν_{eff} through the relation

$$\sigma = \frac{\omega_{pe}^2}{4\pi\nu_{eff}} \quad \ldots (5)$$

The first two lengths in Table II are relevant to a purely 'classical' model (i.e. only binary collisions considered): the mean free path for ion-ion collisions λ_{ii} is clearly too long to account for the shock dissipation, since it is a significant fraction of the total scale length $L \approx 10^{10}$ cm. If $\delta \approx$ ion gyro-radius ρ_L, the current density is large enough to lead to electrostatic instability, thereby invalidating the assumption that only binary collisions need to be considered. The value c/ω_{pi} is the usual scale length observed in collisionless shocks, based on an effective collision frequency $\approx \omega_{pi}$, and c/ω_{pe} is the collision free skin depth, the smallest scale over which, because of electron inertia, we can expect a large magnetic field gradient. Laboratory generated collisionless shocks (with $M > 1$) are observed to have widths $\delta \approx c/\omega_{pi}$. The length u/ω_{pi} corresponds to the initial depth of penetration into a turbulent plasma of the inflowing ions, whether the turbulence is set up as a result of a current driven instability or by ion-ion counterstreaming (Vedenov et al., 1961), and is a measure of the effective distance within which the initial fluid flow can be changed. The corresponding values of j and v_d show that the magnetic field cannot change on the same scale without leading to instability. Finally, the last two lengths are characteristic of magnetic diffusion, $\delta = c^2/4\pi\sigma u$, taking in turn classical conductivity, which, as in the neutral point region, leads to absurdity, and the experimentally determined value $\sigma_T = \frac{1}{2}\left(\frac{m_p}{m_e}\right)^{1/3} \omega_{pe}$ (Hamberger and Friedman, 1968) for highly turbulent plasma, which leads to δ values of the same order as seen in fast shocks. Table II shows that for all scale lengths listed (other than the unlikely $\delta = \lambda_{ii}$) the drift velocity is adequate to satisfy the instability condition, equations (2a) and (2b).

Within the shock front there is a resistive electric field j/σ which we can find from equations (1) and (5). If we take for example the largest physically reasonable value, $\delta = c/\omega_{pi}$, corresponding to $\nu_{eff} \approx \omega_{pi}$

$$E = j/\sigma = \frac{B}{m_p/m_e} \approx 75 \text{ V cm}^{-1} \qquad \dots (6)$$

This is probably an underestimate of the resistive electric field: for comparison it is one half of the total Lorentz electric field available

$$\underset{\sim}{E}_L = -\frac{\underset{\sim}{u} \times \underset{\sim}{B}}{c} \approx 150 \text{ Vcm}^{-1} \qquad \dots (7)$$

which is the resistive electric field in the diffusion region IV. (Friedman and Hamberger, 1968).

In an earlier attempt to find a physically meaningful scale length for the boundary layer Sturrock (1968) assumed it to be turbulent so that conditions are determined by anomalous rather than classical diffusion; by analogy with certain results from controlled fusion experiments, he suggested substituting for the magnetic diffusion coefficient

$$D_M = c^2/4\pi\sigma \qquad \dots (8)$$

the value given by Bohm et al. (1949)

$$D_B = ckT_e/16eB \qquad \dots (9)$$

With his chosen parameters he obtains $\delta \approx 10^2$ cm, of the same order as we suggest.

The notion that the active part of the flare is highly microturbulent is supported by the observation that strong microwave emission occurs at the same time as the electrons become heated, as evidenced by the onset of the X-ray emission. In a later section a more direct comparison with experiment is made based on this assumption.

3. ACCELERATION OF PARTICLES IN A FLARE

A particle moving with velocity $\underset{\sim}{v}$ in the configuration of Fig.(1) will experience a total electric field

$$\underset{\sim}{E}_L = (\underset{\sim}{E})_z + \frac{\underset{\sim}{v} \times \underset{\sim}{B}}{c} \qquad \dots (10)$$

where $(\underset{\sim}{E})_z$ is the constant electric field in the stationary frame

$$(E_z) = -\frac{u_{x_1} B_{y_1}}{c} \qquad \dots (11)$$

If it has a trajectory such that, in moving initially against the flow, it is reflected back through the plane $x = 0$, passes into the opposite flow region, and is again reflected back through $x = 0$, etc., it will gain energy provided there is a net drift of its guiding centre in the z-direction.

Within the highly resistive layers where current is flowing, e.g. in the diffusion region, the electric force is balanced by the frictional force provided by the ensemblage of the fluctuating electric fields, so that the average particle motion is a steady drift between electrons and ions with a relative velocity

$$v_d = \frac{E_z e}{m_e \nu_{eff}} = \frac{\sigma E_z}{ne} \qquad \ldots (12)$$

We are not concerned in this section with the motion of the bulk of the particles, but with a relatively small fraction that will fulfill certain conditions, mentioned later.

Since the magnetic field configuration is very complicated, especially near the neutral line and in the shock fronts where large field gradients occur, and the effect of the fluctuating electric fields is very difficult to include in a calculation, we need to resort to a numerical method to integrate the particle orbits, making certain simplifying assumptions. These are:

(1) we assume the shape of the spatial variation of magnetic field;
(2) we consider only particles which have sufficient energy to escape from the hot plasma region so that they can partake in the Fermi-type acceleration mentioned above.

The implications of these assumptions will be dealt with in separate paper (Friedman, 1969). Here we shall state simply that the first assumption does not restrict the process; acceleration occurs for any assumed form which is compatible with the wave model, while the second condition can be fulfilled provided the particle has sufficient momentum to pass through the turbulent region without its orbit being seriously affected by scattering, i.e. it has an effective mean free path

$$\lambda_{turb} > \delta^* \qquad \ldots (13)$$

Condition (13) will be satisfied, even for a high level of turbulence, by protons (but not electrons) with an initial energy $\varepsilon_1 \gtrsim 5kT \approx 50$ keV in our example, so that for computational purposes we can ignore the effect of the turbulence on the particle trajectories.

We must therefore postulate that the accelerated particles originate from the high energy tail of the distribution of protons in the hot plasma, $T \approx 10^8$K, so that we restrict our attention to relatively few particles, consistent with the statement above.

The actual computation consists of integrating the momentum equation

$$\frac{d}{dt}(m_p \underset{\sim}{v}) = e\underset{\sim}{E}_L \qquad \ldots (14)$$

The particles are taken to have initial velocity corresponding to $\varepsilon_1 = 50$ keV, and randomly chosen directions and positions within the diffusion region IV.

The results (summarized only here) show that about 70% of the particles chosen are accelerated to energy > 0·1 MeV before they cease to gain energy (i.e. they no longer cross the x = 0 plane). The final energy ε is found to satisfy distribution laws of two distinct forms, depending only on the assumed form for $\underset{\sim}{B}$. These are

$$N(>\varepsilon) \sim \exp(-\varepsilon/\varepsilon_0) \qquad \text{... (15)}$$

and

$$N(>\varepsilon) \sim \varepsilon^{-\beta} \qquad \text{... (16)}$$

where $N(>\varepsilon)$ is the number of particles with final energy $>\varepsilon$, and ε_0 and β are constants which depend only on the model parameters, M_0, n and B.

Observations of the solar proton flux produce distributions of both these forms.

In the more detailed presentation it is shown that

$$N(> 0{\cdot}1 \text{ MeV}) \approx 10^{35} - 10^{36}$$

Using equation (16) with a computed constant $\beta \approx 2$ we find

$$N(> 10 \text{ MeV}) \approx 10^{31} - 10^{32}$$

which agrees with observation (Krimigis, 1965)

The limiting energy is

$$\varepsilon_{max} \approx eE_z L \approx 10^{12} \text{ eV} \qquad \text{... (17)}$$

Because the actual numbers depend on the chosen spatial variation of magnetic field, they can be at best only very approximate, but the qualitative agreement with observation is reasonable.

4. COMMENTS ON SOME OBSERVED EMISSIONS FROM SOLAR FLARES

(a) Microwave and X-ray emission

The early stages of most intense flares are accompanied by a burst of microwave emission, which is unpolarized, covers a wide spectrum in the centimetre wavelength range, and whose intensity increases very rapidly, reaching a maximum within about 1 minute. Peak received fluxes are of order 10^{-19} W m^{-2} Hz^{-1}. The duration is typically a few minutes (Smith and Smith, 1963; Kundu, 1965).

There is a high observational probability that events which exhibit such radiation also emit X-rays, which are usually observed in the 20-50 keV range (Winckler, 1963). Their intensity also increases rapidly, and they persist considerably longer (order 15 min).

Recent work (Arnoldy et al., 1968) has shown the correlation extends to the intensities and, for a given event, to the times of onset and of peak emission.

In many laboratory experiments of the type mentioned (e.g. Fanchenko et al., 1964; Hamberger et al., 1967, 1968; Demidov et al., 1968; Lin and Skoryupin, 1968) similar phenomena are exhibited when the plasma is made highly turbulent by applying a very large electric field, i.e. while it exhibits the greatly reduced conductivity according to equation (5). Fig.2(b) shows typical oscillograms taken in the authors' laboratory which may be compared with Fig.2(a), which refers to the flare of March 30, 1966. The time scales differ, of course, since the turbulence is in one case produced by a transient applied electric field and in the flare determined by the duration of the flow, if our hypothesis is correct. The X-ray duration is a measure of the time the heated electrons remain in each system.

In the laboratory we observe an emission of order 10^{-14} W cm^{-3} Hz^{-1} at 3 cm wavelength when the plasma density $n \approx 5 \times 10^{11}$cm^{-3}, electron temperature $T_e \approx 10^8$K. If we tentatively assume similar conditions of turbulence exist in the shock wave region of a typical flare, then we may make a rough estimate (ignoring absorption, scattering etc.) of the received microwave flux expected from a flare of the radiating mechanism in the same. Taking a radiating volume $L_1^2\ \delta^* \approx 10^{20}$ cm ($L_1 < L$) we obtain a peak received flux $\approx 10^{-18}$W m^{-2}Hz^{-1}, which compares reasonably well with observations.

The most likely mechanism for generating the electromagnetic radiation is that based on the non-linear interaction between the plasma oscillations, as first suggested by Sturrock (1961). Since the spectrum consists of plasma waves of frequencies $\omega \ll \omega_{pe}$ and $\omega \approx \omega_{pe}$, individual interactions would give rise to combination frequencies around both $2\omega_{pe}$ and ω_{pe}. However in a large volume with a high energy density of waves we would expect multiple scattering to occur and to lead to a much wider spectrum, as in fact is observed. More detailed measurements of the microwave spectrum from individual flare events might, however, show maxima at ω_{pe} and $2\omega_{pe}$ (similar to the type III radiation) as has been observed in laboratory turbulence (Demidov et al., 1968).

(b) Type V Continuum Radiation

There seems general agreement that the polarized continuum emissions associated with flares are due to synchrotron radiation from charged particles spiralling in the nearby magnetic field. Observations suggest that there is a close correlation between the occurrence of Type V, Type III and solar cosmic rays (protons) (Smith and Smith, 1963). It has been pointed out by Sturrock (1963) that a beam of fast electrons will be rapidly attenuated by plasma instabilities; they would therefore have far too short a range to account for their presence in the regions which are observed to emit Type V radiation. These objections, however, do not apply to a proton flux (Tsytovitch, 1966).

Since we know that a proton flux does exist, and since, as we have suggested, the model may provide a suitable mechanism for generating energetic protons, let us therefore estimate the requirements if protons are to account for the observed emission.

The maximum intensity of synchrotron radiation occurs at a frequency

$$\nu^* \approx \gamma_p^2 \frac{eB}{2\pi m_{po} c} \qquad \ldots (18)$$

Each proton radiates a power

$$P_{\nu^*} = \frac{2e^4\ B^2\ \gamma_p^2}{3\ m_{po}^2\ c^3 \nu^*} \qquad \ldots (19)$$

where

$$\gamma_p = m_p/m_{po}$$

$$m_p = \text{proton mass}$$

$$m_{po} = \text{proton rest mass}$$

We assume

(a) the radiation is from a volume of radius R

(b) no self-absorption

(c) no collective effects, i.e. radiation ~ total number of protons.

(d) mono-energetic protons, with number density n_p.

The flux received at a distance D is

$$S_{\nu*} = 1/3\ (R^3/D^2)\ n_p\ P_{\nu*} \qquad \ldots (20)$$

Let us take for example

$$B = 500\ \text{G}$$

$$\nu* \approx 100\ \text{MHz}$$

$$S_{\nu*} \approx 10^{-20}\text{W m}^{-2}\text{Hz}^{-1}$$

$$= 10^{-17}\ \text{c.g.s.u.}$$

Then from (18) and (19) we require

$$\gamma_p \approx 12\ (\text{i.e. } \varepsilon \approx 11\ \text{GeV})$$

and

$$n_p\ R^3 \approx 10^{30}$$

Observations suggest $R \approx 2 \times 10^{10}$cm (Smith and Smith, 1963) so that the proton density necessary is $n_p \approx 10^{-1}\text{cm}^{-3}$.

A simple calculation shows the optical depth $\tau \gg 1$ with these parameters, so we are justified in neglecting self-absorption. If collective effects are important (i.e. the protons are bunched) the required number of protons is considerably smaller.

Since radiation can escape only when $2\pi\nu* > \omega_{pe}$, the protons must emerge from chromospheric density regions to the corona. The collisional mean free path for these energies is determined by nuclear interactions: even in the chromosphere ($n_p \approx 2 \times 10^{11}\text{cm}^{-3}$) this is of order 10^{14} cm, which is much greater than the required distances. For comparison, the mean free path for radiative cooling $\approx 10^{26}$ cm.

The proton radiation is also subject to synchrotron absorption by electrons surrounding the flare. However, this will be significant only for frequencies near to or above the electron gyro-frequency, $eB/2\pi\gamma_e m_{eo}c$. Equating this to 100 mHz, we find $\gamma_e \approx 14$. Since neither particle measurements nor X-ray observations indicate the presence of significant numbers of electrons with such energy (6·5 MeV) this effect should be unimportant.

It is interesting to note that Kaplan and Tsytovitch (1968) have recently shown that Type III emission can be explained as originating from plasma instabilities caused by a flux of high energy protons (velocity $\lesssim 0{\cdot}9c$) traversing the corona: these can travel the necessary distances (> 10^{10}cm) without serious energy loss, compared with electrons of similar velocities which have an estimated range $\approx 10^5$ cm. The total number density of such protons to account for the observed intensities was estimated to be of order 10 cm^{-3}. Recent laboratory work (Hermann and Fessenden, 1968) has confirmed that a proton beam in a plasma does generate electromagnetic emission at ω_{pe} and $2\omega_{pe}$, similar to that observed from Type III radiation.

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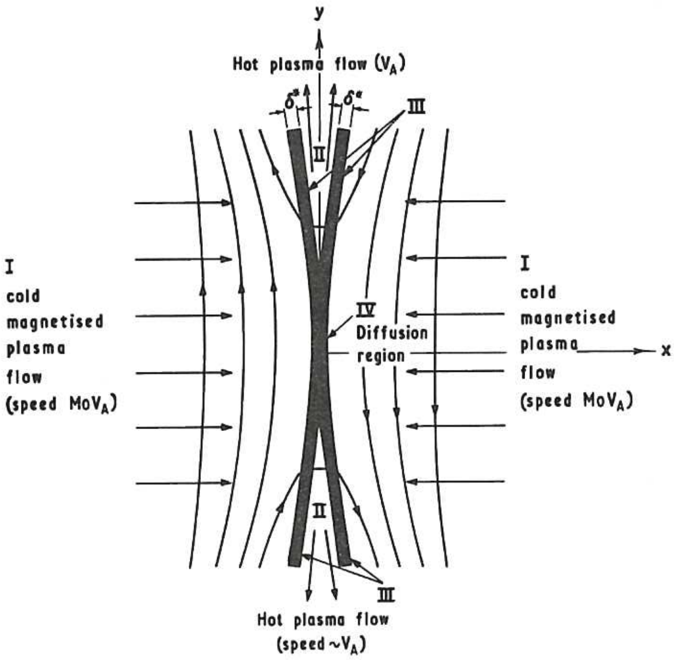

Fig. 1 (CLM-P 191)

Magnetic field and plasma flow situation in the Petschek wave model

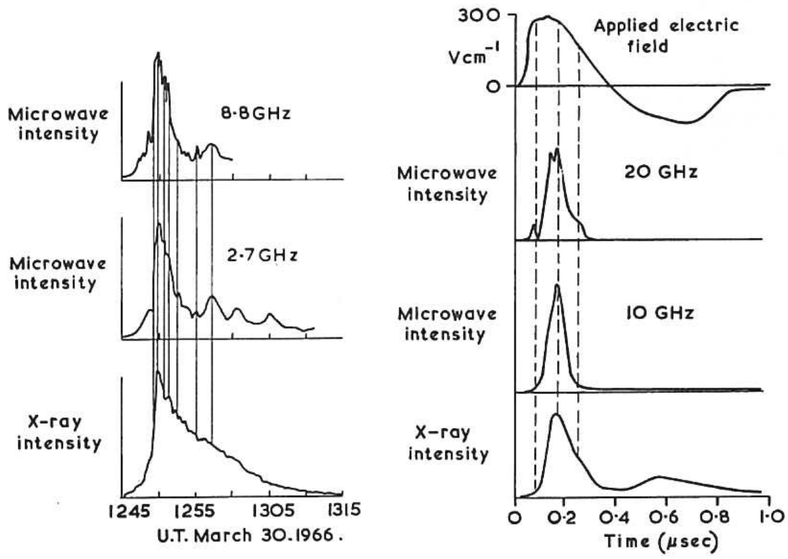

Fig. 2 (CLM-P 191)

(a) X-ray and microwave emission from a typical solar flare (after Amoldy et al., 1968)

(b) X-ray and microwave emission from a laboratory turbulent plasma