MECHANISMS FOR FLASH PHASE PHENOMENA IN SOLAR FLARES

DEAN F. SMITH

High Altitude Observatory, National Center for Atmospheric Research, Boulder, Colo., U.S.A.*

Abstract. Mechanisms for explaining the various forms of particles and radiation observed during the flash phase of solar flares are reviewed under the working hypothesis that the flash phase is the time in which electrons and to a lesser degree protons are accelerated in less than one second. A succession of such accelerations is allowed to explain longer lasting or quasi-periodic phenomena. Mechanisms capable of such acceleration are reviewed and it is concluded that first-order Fermi acceleration in a reconnecting current sheet is the most likely basic process. Such acceleration, however, gives rise to a rather narrow distribution of particle velocities along a given field line which is unstable to the production of electron plasma and ion-acoustic waves. This plasma turbulence can heat the plasma to produce soft X-rays and filter the initially narrow velocity distribution to produce a power law energy distribution. Electrons travelling inward from the acceleration region produce hard X-rays by bremsstrahlung and microwave bursts by gyro-synchrotron emission. Whereas the interpretation of X-ray spectra is relatively straightforward, the interpretation of microwave spectra is difficult because the source at low frequencies can be made optically thick by several different mechanisms.

Electrons travelling further inward presumably thermalize and produce impulsive EUV and Hα emission. The theory for these emissions, although amenable to present techniques in radiative transfer, has not been worked out. Electrons travelling outward give rise to type III radio bursts by excitation of electron plasma waves and the electrons observed at the Earth. Study of the interaction of a stream of electrons with the ambient plasma shows that the electron spectra observed at the Earth do not necessarily reflect their spectrum at the acceleration region since they interact via plasma waves as well as through Coulomb collisions. The mechanisms for the conversion of plasma waves into radiation and the propagation of the radiation from its source to the observer are reviewed.

1. Introduction

We begin by defining the flash phase of a solar flare as that phase in which rapid acceleration of electrons and to a lesser extent protons and heavier nuclei occurs along with a rapid heating of part of the flare plasma. This heating may be due to collisional losses of the accelerated particles, to some part of the acceleration process or to a combination of these two processes. By rapid acceleration of electrons we have in mind a time scale of less than one second. Although it cannot be said that definitive theories exist for the conversion of low energy (1–500 keV) electron energy into all its other manifest forms, it can be said that for all impulsive phenomena which have been examined in detail, the low energy electrons do have sufficient energy. Thus a definition of the flash phase which has rapid acceleration as its basic ingredient should serve as a good working hypothesis on which to try to build a comprehensive theory. While it is clear that any theory for the flash phase must be consistent with the preflare buildup and subsequent flare phases (e.g. act as a trigger for subsequent phases), we limit ourselves to the flash phase as an entity in itself in keeping with the observations which have been presented.

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The basic problem of acceleration of electrons and other particles on the very short (<1 s) time scales required for flash phase phenomena has received little attention in recent years. It has been shown (Smith and Priest, 1972) that the evacuation mechanisms of Alfvén and Carlqvist (1967), Carlqvist (1969), and Syrovatsky (1969, 1972) suffer from basic inconsistencies. These mechanisms were appealing for their simplicity. Takakura (1971) has shown how an electric field can be derived self-consistently by the rotation of a spherical cloud in a magnetic field – a cosmic dynamo. To generate significant electric fields by this mechanism quite large rotational velocities are required. Takakura would overcome this difficulty by placing two such dynamos in magnetic fields of opposite polarity as shown in Figure 1, i.e. in a current sheet, and build up a large current system which would have a large potential drop across it *ala* Alfvén and Carlqvist. This stored energy could be released by allowing the electric field to become large enough to excite electron plasma waves which would reduce the conductivity by several orders of magnitude in the corona. However, Takakura (1971) takes 1 MeV as the maximum energy that a particle can attain which is the potential drop across his current system. This is incorrect in a field of plasma waves since the net acceleration is determined by the microscopic encounters of the particles with the waves rather than the macroscopic potential drop. While there are many attractive features in Takakura's model the configuration of Figure 1 seems too exotic to be a commonly occurring phenomenon on the Sun.

On the other hand, one of the basic ingredients of Figure 1, namely the current sheet, should be a commonly occurring phenomenon and it is well known that large electric fields can be developed when the opposing magnetic fields are reconnecting (Sweet, 1958; Petschek, 1964; Sturrock, 1968). The basic problem here is that, as pointed out by Parker (1973), no one has been able to find a constraint from the reconnection region itself which would determine the maximum rate of reconnection.

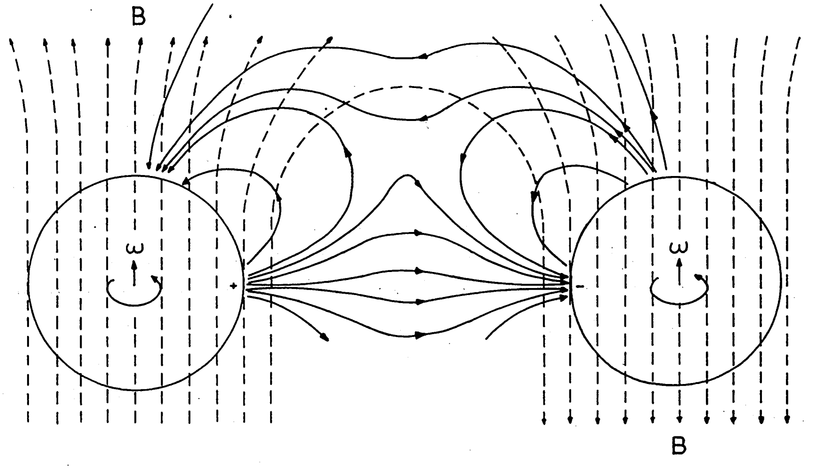

Fig. 1. Schematic picture showing lines of electric force between two spherical clouds rotating in the same direction in magnetic fields of opposite polarities (dashed lines) in Takakura's model.

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Thus Yeh and Axford (1970) have surmised that the rate of reconnection is determined by boundary conditions. It should also be noted that the ideal configurations studied by Petschek (1964), Sonnerup (1970), Yeh and Axford (1970), and Coppi and Friedland (1971) are unlikely to occur in nature since in general there will be a component of magnetic field perpendicular to these two dimensional configurations, as shown in Figure 2, i.e. field lines in general will reconnect at an angle other than 180°. It can be shown by arguments similar to those used by Yeh and Axford (1970) and Parker (1973) for two dimensional configurations that the decrease in the rate of reconnection in a steady state due to this additional component of the field is also undetermined by anything in the reconnection region itself (Cowley, private communication). In the face of such indeterminacy, about all we can do is to look at indicative computer results such as a nonlinear analysis of the tearing mode instability (van Hoven and Cross, 1973), laboratory experiments (Bratenahl and Yeates, 1970) and solid evidence of reconnection in the vicinity of the Earth (Aubry *et al.*, 1970) to affirm our faith that it occurs on the Sun. Having taken that step it is natural to assert that the flash phase is that interval of time in which extremely fast reconnection has begun, but not gone so far as to preclude very coherent phenomena such as rapid acceleration of particles taking place.

Given that the reconnection geometry of Figure 2 is the simplest configuration in which acceleration can take place, one can then study the fate of individual particles in this geometry which has been started by Speiser (1965), Friedman (1969), Cowley

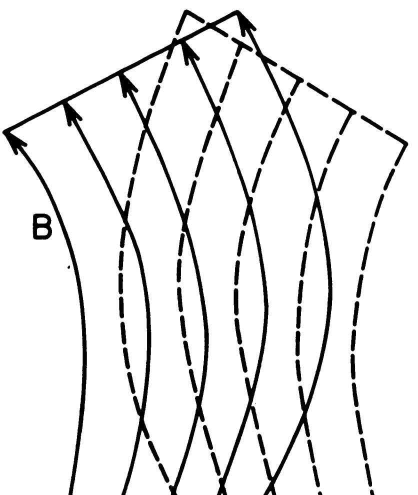

Fig. 2. Schematic drawing of how field lines with significant antiparallel components are likely to occur in nature with field components in every direction.

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(1972), and Eastwood (1972, 1974). We shall review and extend these results in Section 2 and see to what extent they fit the observations.

With nonthermal particles, mostly electrons, the question of how these particles interact to give the myriad forms of radiation observed arises. If we place the height of the primary acceleration region at about 10000 km, then electrons moving down field lines will give rise to microwave bursts by gyro-synchrotron emission and hard X-rays by bremsstrahlung (Holt and Ramaty, 1969; Takakura, 1972), and impulsive EUV by free-bound, bound-bound and bremsstrahlung emission (Kane and Donnelly, 1971) and Hα kernels (Zirin and Tanaka, 1973) by bound-bound emission from excited partially ionized plasma. The theories for these emissions are reviewed in Sections 3 and 4. The electrons travelling up field lines give rise to type III radio bursts by excitation of plasma waves and the electron events observed near the Earth (Smith, 1974a). The theory for type III bursts is reviewed in Section 5. There is no theory for the propagation of low energy electrons in the interplanetary medium, but some facets of this problem are included in Section 5.

Finally there is the heating of the flare plasma at the flash phase which causes the rise of soft X-ray emission that continues into the flare proper (Kahler and Kreplin, 1970; Thomas and Teske, 1971). In the concluding section we check to see how this emission fits in to the theory reviewed and summarize what remains to be done to improve our understanding of the flash phase.

2. Acceleration of Particles in Current Sheets

We concentrate on electrons and recall the requirements provided by the observations. The acceleration mechanism must provide electrons in less than 1 s with a power law distribution.

$$\frac{\mathrm{d}n(E_e)}{\mathrm{d}E_e} \propto E_e^{-\delta} \tag{2.1}$$

with spectral exponent δ in the range 2–5 (Kane, 1973), where $n(E_e)$ is the number of electrons per cm^3 with energies in the range E_e to $E_e + \mathrm{d}E_e$. The spectrum often shows a break at about 100 keV and is harder above 100 keV, especially for large flares (Kane and Anderson, 1970; Frost and Dennis, 1971). For example $\delta = 2.3$ may characterize the spectrum to 100 keV and $\delta = 4.6$ may characterize the spectrum beyond 100 keV. The total number of electrons accelerated above 22 keV is about 10^{36} (Lin and Hudson, 1971) which represents about 10% of the total flare energy (Kane, 1973). In other words the acceleration mechanism must be highly efficient. This requirement does not necessarily rule out stochastic processes, but it does imply that the only wave modes which can be involved are electron plasma waves or radiation since they are the only modes which can exist in the corona in a sufficiently lossless regime. Since the simplest manner of exciting these oscillations is by a stream of suprathermal electrons, the efficiency and time requirements imply that at least the initial stream of electrons must be accelerated directly. In other words electron plasma waves or radiation can at most act as a filter for an already directly accelerated electron stream. Consideration of collision losses (Syrovatsky and Shmeleva, 1972; Cheng, 1972; Biswas and Radhakrishnan, 1973) imply that acceleration will be most effective in the corona which may, however, extend to quite low heights (e.g., 2000 km) in some parts of an active region (Athay, private communication).

We also recall some basic requirements for particle acceleration provided by theory. From the force equation

$$\dot{\mathbf{v}} = \frac{e}{m_e}\left(\mathbf{E} + \frac{1}{c}\,\mathbf{v}\times\mathbf{B}\right), \tag{2.2}$$

where $\mathbf{v}$ is the particle velocity, $\mathbf{E}$ is the electric field and $\mathbf{B}$ is the magnetic field, acceleration requires either an electric field or a timedependent magnetic field. Electric space charge fields due to a charge imbalance via Poisson's equation

$$\nabla\cdot\mathbf{E} = 4\pi e(n_e - n_i), \tag{2.3}$$

where n_e and n_i the electron and ion densities, are eliminated by plasma processes on a time scale ω_{pe}^{-1}, where $\omega_{pe} = (4\pi n_e{}^2/m_e)^{1/2}$ is the plasma frequency. For a density of $3\times10^3\ \mathrm{cm}^{-3}$, the time scale is 4×10^{-10} s which effectively eliminates this possibility in the corona. The only other possibility is an electric field arising from a magnetic field whose lines of force are moving in the system in which acceleration takes place. This induced electric field $-(1/c)\,\mathbf{V}\times\mathbf{B}$, where $\mathbf{V}$ is the velocity of field lines in the frame in which acceleration is taking place, is perpendicular to $\mathbf{B}$ and so leads to an $\mathbf{E}\times\mathbf{B}$ drift across the field lines at a rate

$$\mathbf{v}_E = c\,\frac{\mathbf{E}\times\mathbf{B}}{B^2} = +\mathbf{V}. \tag{2.4}$$

However, once a particle has left the region where the induced electric field exists, which it can do after making half a gyration in the field $\mathbf{B}$, it will have acquired the velocity of the moving magnetic field line as well since the $\mathbf{E}\times\mathbf{B}$ drift is across this field line so that the net velocity gain is $2\mathbf{V}$. In other words, as pointed out by Fermi (1954), the particle effectively makes an elastic collision with the field line. It can be shown (Hayakawa *et al.*, 1964) that all types of magnetic acceleration reduce either to the Fermi mechanism or betatron acceleration (Swann, 1933). The advantage of the Fermi mechanism for our purposes is that it increases the particle velocity along field lines whereas betratron acceleration increases the velocity perpendicular to the field lines. When the velocity of a particle systematically increases, the acceleration process is referred to as a first order process.

We examine particle trajectories in reconnecting current sheets to see to what extent these requirements are satisfied. The simplest two-dimensional configuration devoid of complications such as slow shocks (Sonnerup, 1970) is shown in Figure 3. Field lines start to move in towards $y=0$ from both sides, for reasons which we shall not treat in view of the problems enumerated in the introduction, with a velocity far from the sheet v_f where the z-component of the magnetic field is B_f. As the field lines approach

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the neutral point they bend in towards the neutral point, are reconnected and move out in both directions along z. As shown in Figure 3, the motion of field lines gives rise to an electric field

$$E_x = \frac{v_f B_f}{c}. \tag{2.5}$$

A particle which is injected so that it is travelling exactly along this axis will be accelerated indefinitely by this field as shown by Speiser (1965). It sees an infinite

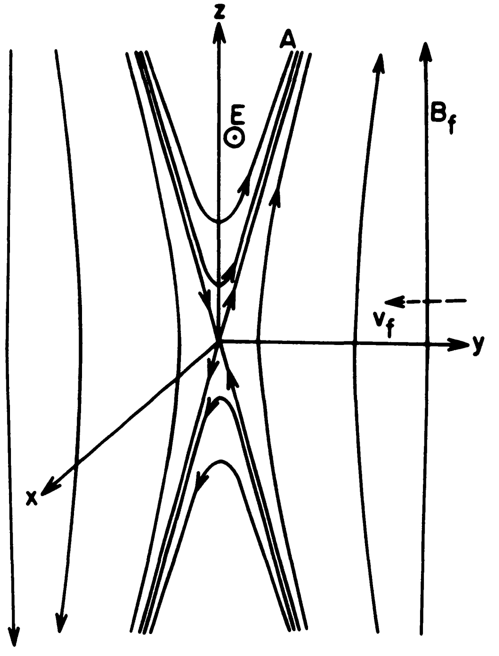

Fig. 3. Qualitative picture of the reconnection process showing the coordinate system employed. Magnetic field lines of strength B_f move toward the neutral point with velocity v_f which gives rise to an electric field E_x.

succession of first order Fermi accelerations and its final energy is limited only by the size of the system. However, the number of particles which satisfy the very stringent condition of being injected right along the neutral line is extremely small. The vast majority of particles entering the current sheet end up where there is a finite transverse field B_y. Thus we consider the fate of these particles in detail.

As first shown by Speiser (1965), one of the simplest ways to treat this problem is to make a transformation to a system moving with a field line since in this frame the electric field vanishes. We consider the field line labelled A in Figure 3. In the Sun's frame of reference, the incoming velocity of an electron is given by

$$\mathbf{v}_0 = \mathbf{v}_{\parallel 0} \pm \mathbf{v}_{c0} + \mathbf{v}_E, \tag{2.6}$$

where $\mathbf{v}_{\parallel 0}$ is the velocity parallel to $\mathbf{B}$ at time $t=0$, $\mathbf{v}_{c0}$ is the initial perpendicular velocity and $\mathbf{v}_E$ is the drift velocity given by Equation (2.4) with $\mathbf{B}$ the field strength of field line A and $\mathbf{E}$ given by Equation (2.5). Transforming to the rest frame of the field line (primed frame) which is moving with speed E_x/B_y in the z-direction, Equation (2.6) becomes

$$\mathbf{v}'_0 = \mathbf{v}'_{\parallel 0} + \mathbf{v}'_{c0}, \tag{2.7}$$

where the transformation velocity has been absorbed into $\mathbf{v}'_{\parallel 0}$. If we neglect the thermal velocity v_e of the electron because we are interested in transformation velocities somewhat higher than v_e, then the electron may simply be assumed to be streaming into the reversal region with a velocity $\mathbf{v}'_0 = \mathbf{v}'_{\parallel 0}$, where $\mathbf{v}'_{\parallel 0} \cdot \hat{z} = v'_{z0} = E_x/B_y$ and $\mathbf{v}'_{\parallel 0} \cdot \hat{y} = v'_{y0} = = E_x/B_z$. Thus the initial electron velocity is specified by E_x, B_y and B_z. Particle motion in this frame, which has been studied both analytically and numerically by Speiser (1965) and Eastwood (1972), consists simply of gyrations about the B_y and B_z fields. A particle oscillates between the reversing B_z field and is gradually turned so that it is moving in the positive z-direction by the B_y field which must be much weaker for v'_{z0} to be large. What is important for our purpose is that after half a gyration in the B_y field, a particle is ejected almost along the field line with a velocity $2v'_{z0}$ in the Sun's frame of reference. In other words, aside from a slightly more complicated motion due to the non-adiabatic nature of a current sheet, a particle is accelerated just as in a first-order Fermi process and because $B_y/B_z \ll 1$, almost along B_z. The energy gained by a particle is directly proportional to its mass as is the time for acceleration

$$\tau_A = \frac{\pi mc}{eB_y}. \tag{2.8}$$

The transverse field B_y is a very important parameter in this process since together with E_x it determines the energy gain

$$\Delta E = 2mv_{z0}^2 = \frac{2mv_f^2 B_f^2}{B_y^2}, \tag{2.9}$$

where we have used Equation (2.5). Some typical values for ΔE for electrons are given in Table I. Friedman (1969) showed that by allowing B_y to vary as $z^3/|z|$, a power law distribution of particle energies would be obtained with the index δ depending on $B_{z0} = B_f$, the plasma density n_e, the Mach number v_f/v_A where v_A is the Alfvén velocity in B_f and the spatial variations of the magnetic field. However, Friedman failed to explain how the different energy particles going along different field lines will get mixed. In fact, as pointed out by Eastwood (1972, 1974), the distribution of particle velocities along a given field line at one time is quite narrow and going in the opposite direction to the incoming particles as shown in Figure 4. This type of distribution is unstable to excitation of electron plasma waves if the distance between peaks is greater than v_e which can easily be satisfied according to Table I and ion-acoustic waves by the two-stream instability (Stringer, 1964). The growth rate for electron plasma waves in this case is about 0.4 ω_{pe} so that 20 e-folding steps take 1.8×10^{-8} s in a plasma of density 3×10^9 cm^{-3}. Thus, if the reconnection region has any reason-

TABLE I

Energy gains of electrons in a reconnecting current sheet

v_f (cm s^{-1})	B_f (G)	B_y (G)	E (keV)
		0.5	1.2
10^6	500	2.0	0.1
		5.0	0.01
		0.5	11.5
3.1×10^6	500	2.0	0.7
		5.0	0.1
		0.5	115
10^7	500	2.0	7.2
		5.0	1.2

able finite thickness like 0.1 km (Sturrock, 1968), then this instability is unavoidable in the reconnection region itself. Friedman (1969) considered such plasma turbulence to be highly likely, but he considered the turbulence to be isotropic. This led him to conclude that electrons could not be efficiently accelerated because they would be scattered and only ions with a quite high injection velocity would be unaffected, i.e. the plasma turbulence would act as a selection mechanism to allow a few ions to reach very high energies. If this indeed were the case, this mechanism would clearly be unsuitable for our purpose.

However, strongly excited electron plasma waves and ion-acoustic waves in a magnetic field which is sufficiently strong so that $\omega_{He} \approx \omega_{pe}$, where $\omega_{He} = eB/m_e c$ is the electron Larmor frequency, are much more nearly one-dimensionally distributed along the magnetic field. The reason for this behavior for electron plasma waves can be seen from the dispersion relation (Kaplan and Tsytovich, 1973)

$$\omega_p(\mathbf{k}) = (\omega_{pe}^2 + \omega_{He}^2 \sin^2\theta + 3v_e^2 k^2)^{1/2}, \qquad (2.10)$$

where θ is the angle between the direction of the magnetic field and the wave vector $\mathbf{k}$. It is a characteristic of nonlinear interactions which must occur when the energy density in plasma waves W_p becomes sufficiently high that the frequency of the plasma

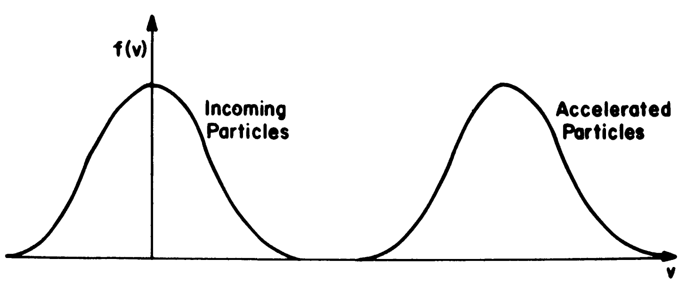

Fig. 4. The distribution of particle velocities along a given field line at one time which consists of the incoming and accelerated particles with their thermal spreads.

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wave decreases somewhat in the process. In contrast to the case in the absence of a magnetic field where this decrease can only occur by a decrease in the wave number k, it can be seen from Equation (2.10) that in a magnetic field the decrease can also occur by a decrease in the angle θ. As a result plasma waves become distributed almost completely along the magnetic field. Ion-acoustic waves in a magnetic field $\omega_{He} \approx \omega_{pe}$ tend to be unaffected by the magnetic field, but for high levels of the energy density in these waves, W_s, become subject to resonance broadening (Tsytovich, 1971). This has the effect of keeping the waves close to the direction in which they were excited which is primarily along the magnetic field since the growth rate maximizes here and is consistent with experiments where the angular distribution of waves has been measured directly (Paul *et al.*, 1972).

As a consequence of this almost one dimensional plasma turbulence electrons will be preferentially accelerated along the field lines. We have in effect overcome the problem of electric space-charge fields by providing a fluctuating field. The nonlinear evolution of a spectrum of electron plasma waves in a field of ion-acoustic waves has not been worked out in detail, but the processes which will be important can be enumerated. The most important process is the induced decay of a plasma wave p into another plasma wave p' and an ion-acoustic wave s (Tsytovich, 1970)

$$p \to p' + s. \tag{2.11}$$

The wave p' has a smaller wave number k' than p since some momentum goes into the ion-acoustic wave, and thus a larger phase velocity $v'_{\text{ph}} = \omega_{pe}/k'$. As a result of many decay processes of type (2.11) coupled with the fact that when the p and s waves have approximately equal wave numbers the process

$$p + s \to p' \tag{2.12}$$

which increases v'_{ph} will complete with process (2.11), a quasi-stationary spectrum of plasma turbulence will be set up (Pikel'ner and Tsytovich, 1968). This spectrum may have the form shown in Figure 5 with $k_0 < \omega_p/c$ so that the only part of the spectrum for which $v_{\text{ph}} < c$ which can interact with particles is a power law of the form

$$W_p(k) \propto k^{-\nu}. \tag{2.13}$$

The type of particle spectrum resulting from such a distribution of waves can be studied by means of the Fokker-Planck equation (Tsytovich, 1970)

$$\frac{\mathrm{d}f_\varepsilon}{\mathrm{d}t} = D\frac{\partial^2 f_\varepsilon}{\partial \varepsilon^2} - 2D\frac{\partial}{\partial \varepsilon}\left(\frac{f_\varepsilon}{\varepsilon}\right), \tag{2.14}$$

where f_ε is the distribution of particle energies ε and D is the diffusion coefficient given by (Tsytovich, 1970)

$$D = \frac{2\pi^2 e^2 \omega_{pe}^2}{v^3} \int_{\omega_{pe}/v}^{\infty} \frac{W_p(k)}{k^3}\,\mathrm{d}k. \tag{2.15}$$

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The first term on the right hand side of Equation (2.14) describes random acceleration while the second term defines systematic acceleration. For the power law spectrum of waves (2.13),

$$D \propto \varepsilon^{2(\nu-1)/2} \tag{2.16}$$

(Pikel'ner and Tsytovich, 1968) and the resulting differential electron energy spectrum

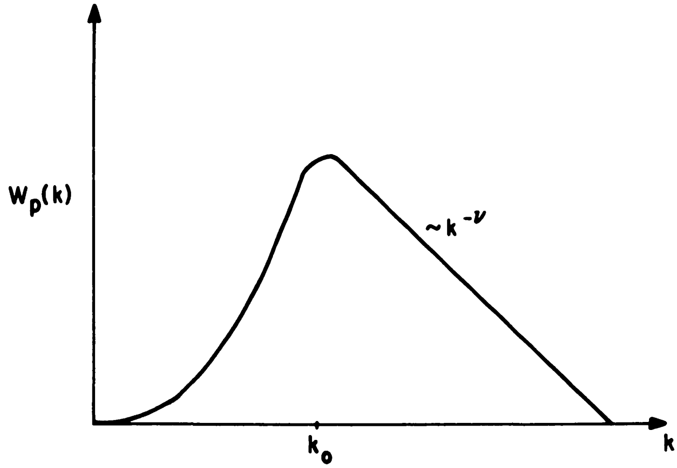

Fig. 5. A schematic of the plasma wave distribution along a field line resulting from strong nonlinear interaction of plasma waves. The wave number k_0 characterizes the beginning of the power law spectrum of waves.

has the form of Equation (2.1) with $\delta = 2$–5 for $\nu = 5$–11. Thus, although it has not been shown in detail how plasma wave spectra of the form of (2.13) with $\nu = 5$–11 are formed, we have shown that such spectra will lead to the observed electron spectrum. Moreover, since δ is determined by microscopic processes such as (2.11) rather than several macroscopic parameters as in the case of Friedman's analysis, the range of δ should be much smaller since microscopic processes tend to lead to well defined spectra (Tsytovich and Chikhachev, 1970). One of the most important effects of the 'plasma wave filter' will be to produce 100 keV electrons along field lines which originally had only 12 keV electrons.

In summary we have found that the initial direct acceleration of particles in current sheets leads to the preferential acceleration of electrons because, although ions obtain m_i/m_e times more energy, it takes them m_i/m_e more time to obtain this energy (Equation (2.8)). The accelerated electrons and ions flowing back into the incoming plasma are unstable which leads to production of electron plasma waves and ion-acoustic waves which are almost aligned along the field. Nonlinear processes lead to a spectrum of plasma waves which leads to a power law distribution of electrons. Since there are many more electrons which can resonate with the plasma waves, electrons again will be preferentially accelerated in contrast to the conclusion reached by Friedman (1969). Since the growth time for plasma waves and nonlinear processes is at most a few micro-seconds, it is readily possible to satisfy the acceleration time requirement. There is no way without a detailed analysis to determine the efficiency of such an acceleration process. However, the initial direct acceleration in the current sheet converts magnetic energy into particle kinetic energy at close to unit efficiency. The electron plasma waves are rapidly transferred to high phase velocities and thus subject only to collision damping. The ion-acoustic waves, on the other hand, are subject to very strong damping until the electron temperature T_e becomes much higher than the ion temperature T_i by damping of both wave modes. Thus an efficiency of 10% would certainly seem to be within the capabilities of this mechanism, but it remains to be shown. The observation of a small number of 50 MeV protons at the time of a type III burst (MacDonald and van Hollebeke, 1973) is consistent with this mechanism since some of the protons directly accelerated by the current sheet will be further accelerated by the plasma waves.

The increase in the electron temperature which occurs on a time scale which is long compared to the acceleration time will eventually raise the threshold for instability for plasma waves and increase the relative importance of the last term in Equation (2.10). In other words the plasma waves which are excited will become more three-dimensional and the situation envisaged by Friedman (1969) will take place where the electrons which are directly accelerated are scattered by the plasma and ion-acoustic waves and effective acceleration ceases except for a very select group of ions. Thus, while it cannot be said that we have a quantitative theory for efficient electron acceleration at the present time, the basic components of such a theory seem to be well in hand. The triggering mechanism for this process is the triggering mechanism for magnetic field reconnection and remains an open question until the conditions determining the rate of reconnection become clear. The addition of a third component of the field to Figure 3 as in Figure 2 will not significantly alter the results presented as long as neutralization of accelerated particles can occur sufficiently quickly to rapidly balance any charge imbalance. We shall defer a discussion of the height of the acceleration region to Section 5.

3. Impulsive Microwave and Hard X-Ray Emission

The nonthermal electrons travelling down field lines in a region close to or contiguous with the acceleration region give rise to microwave and hard X-ray bursts. Because each feature in the microwave intensity profile is followed by the X-ray intensity profile, most theories have attempted to explain both emissions from sources which are in total coincident (Holt and Ramaty, 1969; Takakura, 1972). Both of the cited analyses assumed that X-rays were produced by a thin-target bremsstrahlung process. The thin vs thick-target controversy is now the principle problem in interpreting X-ray spectra. The rapid rise time (~ 1 s), power law spectrum and tracking of impulsive microwave bursts which must be produced by nonthermal electrons (Ramaty and Petrosian, 1972) argue very strongly for the nonthermal nature of hard ($\gtrsim 20$ keV) X-rays. In the thick-target model electrons lose all of their energy through

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collisions in the X-ray emitting region so that the spectrum is a result of the equilibrium between the injection of newly accelerated electrons and the loss of electrons through collisions. In the thin-target model, electrons escape from the X-ray region in a time short compared to their collisional lifetime, and the spectrum of injected and X-ray emitting electrons is the same since collisions are negligible.

The main proponents of a the thick-target model are Brown (1972) and Hudson (1972). While their arguments are valid for large flares, they may meet with serious difficulty for the majority of flares which are small because:

(1) There are no major differences in disk flare spectra and one behind-the-limb flare spectrum which was produced high enough so that it was probably due to thin-target emission (Kane, private communication) as would be expected for thick-target emission from disk flares since γ in the relation

$$\frac{\mathrm{d}J_{h\nu}}{\mathrm{d}(h\nu)} = A(h\nu)^{-\gamma}\ \mathrm{cm}^{-2}\ \mathrm{s}^{-1}\ \mathrm{keV}^{-1} \tag{3.1}$$

decreases by a factor of about 2 for thick target emission.

(2) The electron spectra observed at the earth are consistent with thin-target emission for several measured X-ray events on a statistical basis (Kane, 1973) and for one small flare where the spectra were measured concurrently (Datlowe and Lin, 1973).

Thus we can say that the possibility of primarily thick-target emission in most flares is not favored by simple interpretations of data and concentrate on thin-target emission. Even within this realm there is the question of whether the injection of electrons into the X-ray source is impulsive or continuous. In the impulsive model (Takakura and Kai, 1966) electron injection begins at the onset of the X-ray burst and stops at the time of X-ray maximum. In the continuous injection model (Kane and Anderson, 1970; Syrovatskii and Shmeleva, 1972; Brown, 1972) the electron acceleration process determines primarily the time intensity profile of the entire X-ray burst, the time constants for trapping and/or collision loss being assumed short in comparison to the characteristic time for acceleration. At photon energies greater than about 30 keV, the continuous injection model is favored because the decay of the burst often has a profile much like the rise of the burst (Kane and Anderson, 1970), but at lower energies the spectra may be contaminated by a thermal component and it is difficult to provide evidence to support either hypothesis. Because the continuous injection model is favored at higher energies, we shall assume that it is applicable to the whole nonthermal component of interest here.

In this range X-ray production is simply due to nonthermal bremsstrahlung and leads to a spectrum of X-rays at the Earth of (Kane and Anderson, 1970)

$$\frac{\mathrm{d}J(E)}{\mathrm{d}E} \approx 5.6 \times 10^{-52}\,\frac{n_i V}{E} \int_{E}^{100} \frac{1}{E_e}\,\frac{\mathrm{d}J_e(E_e)}{\mathrm{d}E_e} \times$$

$$\times \ln\left[\left(\frac{E_e}{E}\right)^{1/2}+\left(\frac{E_e}{E}-1\right)^{1/2}\right] \mathrm{d}E_e \text{ photons cm}^{-2}\text{ s}^{-1}\text{ keV}^{-1}, \quad (3.2)$$

where $E=h\nu$ is the photon energy, n_i is the ion density and $\mathrm{d}J_e(E_e)/\mathrm{d}E_e$ is the electron spectrum which is of the form of (2.1), but with exponent $\delta'=\delta-\frac{1}{2}$. In particular, for this form of $\mathrm{d}J_e(E_e)/\mathrm{d}E_e$, the X-ray spectrum (3.2) will have the form of a power law. Although this type of analysis is a gross simplification because the anisotropy in pitch angles, inhomogeneous density structure of the source, etc. are neglected, we are confronted with the following situation. All more sophisticated analyses such as those of Brown (1972) which take into account the anisotropy and change of pitch angles in the source must do so by introducing additional parameters. These parameters taken singly inevitably have the undesirable effect of requiring a change in the electron spectral index δ'. On the other hand the possible virture of Equation (3.2) is that it predicts spectra consistent with the range of electron spectra observed at earth on a statistical basis (Kane, 1973) and on a one-one basis for one event (Datlowe and Lin, 1973). Thus, while Equation (3.2) ignores a considerable amount of basic physics, any more sophisticated approach must introduce at least two additional parameters which, at least in small flares, have opposite effects on the required electron spectrum, an approach which hardly seems better with present observations unless it can be shown that use of Equation (3.2) leads to an incorrect estimate of the number of electrons required. This may well be the case in large flares as the analysis of Brown (1972) would indicate. In the case that thick target processes are dominant, an even simpler equation than (3.2) results (Brown (1972)). The only case of difficulty is the intermediate target case where both thin and thick target processes play a role. Since thick target emission is several hundred times more efficient than thin target emission, it is likely to be dominant when present and the intermediate target case may not be a problem.

In contrast to the simplified approach useful for most X-rays, a quite involved approach is required to make the microwave and X-ray sources coincident in total (Holt and Ramaty, 1969; Takakura, 1972) and to explain the rather large variability of microwave spectra (Ramaty and Petrosian, 1972). The source of this difficulty lies largely in the fact that the microwave source is optically thick in some region whereas the X-ray source can always be considered optically thin. Since there are several possible absorption mechanisms one is forced to consider all the basic physics including the inhomogeneity of the magnetic field. Thus we limit our review to a qualitative discussion of which effects seem to be important. The emission mechanism is gyro-synchrotron emission (Ramaty, 1969). The most important absorption mechanism operative at low frequencies (1–20 GHz) in most sources is self-absorption of the emitted radiation by the emitting electrons (Holt and Ramaty, 1969; Takakura, 1972). The frequency of maximum emission ν_{sa} resulting from self-absorption of a relativistic synchrotron source was given by Slish (1963)

$$\nu_{sa}\approx 10^{12.5}(S_m/\Omega)^{2/5}\, B_{\perp}^{1/5} \quad (3.3)$$

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and was shown to agree with numerical results (Ramaty, 1969) to within 10% over the range of interest for most microwave bursts (Ramaty and Petrosian, 1972). Here S_m/Ω is the maximum brightness and $B_\perp$ is the magnetic field component perpendicular to the line of sight in the emitting region. Both Holt and Ramaty (1969) and Takakura (1972) have shown that Razin-Tsytovich suppression is negligible for densities less than or of order 10^{10} cm^{-3}. Another absorption mechanism which is important for the relatively rare microwave bursts with flat spectra is free-free absorption (Zirin *et al.*, 1971) which leads to

$$\nu_{ff} = 0.39(EM)^{1/2}\, T_e^{-3/4}\, R^{-1}\, \Omega^{-1/2} \tag{3.4}$$

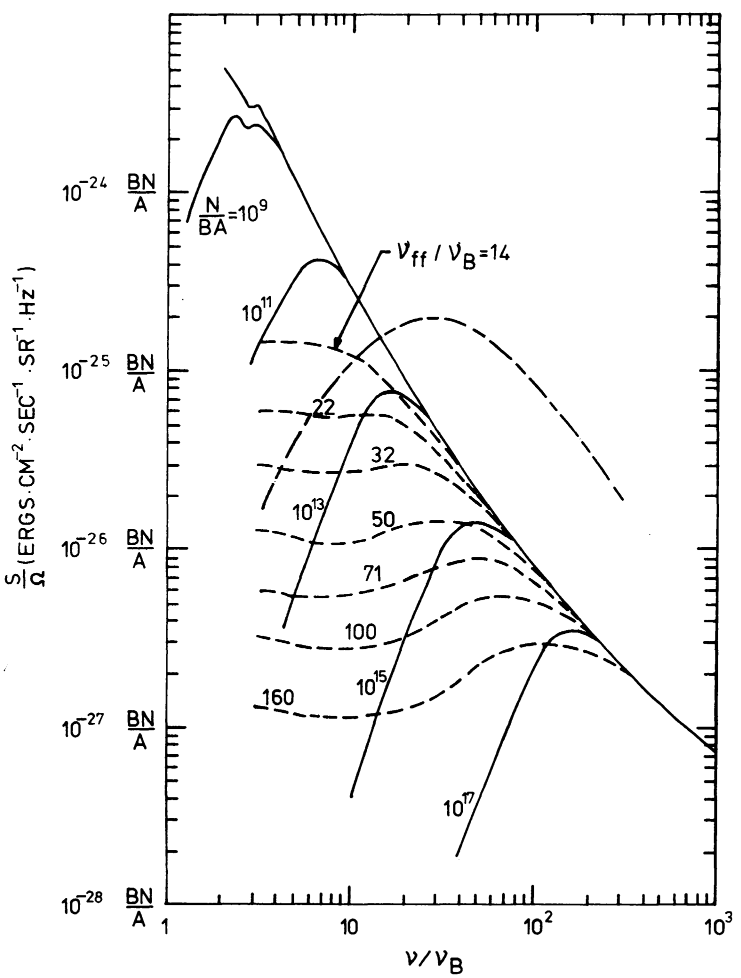

Fig. 6. Gyrosynchrotron spectra from a nonthermal electron distribution. Ramaty's and Ramaty and Petrosian's results for a source with a uniform magnetic field: *Heavy solid line*, gyrosynchrotron emissivity; *light solid line*, self-absorbed brightness; *dashed line*, free-free absorbed brightness. Takakura's result for self-absorbed brightness in a non-uniform magnetic field, *long-short dashed line*, showing the broadening effect of nonuniformity.

for the frequency corresponding to unit optical depth (Ramaty and Petrosian, 1972), where $EM = n_e^2 V$ is the thermal emission measure, R is the distance to the source of volume V and Ω is the solid angle which it subtends. Another possibility is absorption below the plasma frequency (Zirin *et al.*, 1971).

The effect of a non-uniform magnetic field, in particular a dipole field, together with self-absorption was studied by Takakura (1972). His results for a dipole of strength 2500 G are shown in Figure 6 for a height at which $B = 60$ G together with those of Ramaty and Petrosian for free-free and self-absorption in a uniform field. We see that the effect of a nonuniform magnetic field is to broaden the spectrum and make it look more like many observed spectra. Each part of the spectrum comes from a limited part of the source which decreases the apparent problem of too many electrons for the observed fluxes found by earlier workers (Peterson and Winkler, 1959; Holt and Cline, 1968) for coincident X-ray and microwave sources. Anisotropy of electrons with most of the energy along the field also acts in this direction as does a high energy cutoff (Holt and Ramaty, 1969). In short whereas X-ray production seems comparatively insensitive to source parameters in the thin-target regime, microwave production is very sensitive to these parameters and there are several possibilities for decreasing the microwave production from a given electron population. This in turn has the undesirable effect of making interpretations of microwave spectra extremely model dependent so that until other parameters in active regions become better known, they offer little possibility of increasing our knowledge of these regions. In particular, selection of one cutoff mechanism over another with qualitative arguments (Zirin *et al.*, 1971) seems dubious.

From the analysis of impulsive hard X-ray and microwave bursts (Holt and Ramaty, 1969; Takakura, 1972), we can place the following limitations on their source. The electron density $n_e \lesssim 10^{10}$ cm^{-3} and the magnetic field $B \lesssim 350$ G. The total number of electrons with energies greater than 100 keV, $N(>100 \text{ keV}) \lesssim 10^{37}$ and these electrons have a power-law spectrum with exponent $\delta' = 3-5$. The total amount of energy involved is up to 1.6×10^{30} erg, but is more typically a few times 10^{28} erg in interpretations of a small solar flare (e.g., Kane, 1973).

4. Impulsive EUV and Hα Emission

As some of the nonthermal electrons descend into the chromosphere, they excite EUV radiation (Kane and Donnelly, 1971; Donnelly *et al.*, 1973) and Hα (Zirin *et al.*, 1971; Vorpahl, 1973). Since the highest energy electrons which are subject to least scattering will descend the deepest (Brown, 1972; Hudson, 1972), we are confronted with the fact that these electrons must also produce X-rays by thick-target bremsstrahlung processes. From the results of Section 3 it can be argued that the percentage of flares in which thick-target bremsstrahlung is the dominant source of X-rays is probably small so that only a small fraction of the accelerated electrons penetrate deep into the chromosphere in most flares. The electrons which do manage to make their way are rapidly thermalized in the chromosphere due to Coulomb and ionizing collisions (Brown, 1973) which results in a much larger number of quasithermal electrons with higher effective temperatures than normal. These electrons in turn excite atoms as in a thermal plasma to produce EUV and Hα emission.

The only attempt at any quantitative calculation of the impulsive EUV or Hα emission has been made by Kane and Donnelly (1971) for the continuum part of the spectrum produced by recombination assuming the line contribution is small. Since short of 512 Å there is always an enhancement of chromospheric line emission during the flash phase (Hall and Hinteregger, 1969) whereas there is no firm evidence of any continuum enhancement (Neupert, private communication), complete neglect of the line component is unjustified (see also Kane, 1974). Thus we do not reproduce Kane and Donnelly's computation. Rather, following the approach of Brown (1973), we indicate what needs to be done. Brown derived the temperature and density profiles for a model atmosphere in quasi-equilibrium with a power law spectrum of electrons injected vertically downward treating the dominant Balmer continuum correctly in contrast to Hudson (1972) who neglected it by use of a one-level atom. However, he still treated the radiative losses incorrectly which requires use of a three-level atom (Canfield, private communication). Since a part of the chromosphere is 'boiled off' in this approach in which constant pressure is assumed, a hydrodynamic problem must be solved, but the radiative transfer can be treated as a steady-state problem.

In the case of interest for most impulsive phenomena in flares, on the other hand, the time scale for gas motion t_D which is given roughly by the time for a sound wave to traverse one density scale height H is larger than the time scale for electron injection t_B. For example, for $T=10^4$ K and $H=750$ km, $t_D \approx 50$ s whereas most impulsive injection occurs on a time scale less than t_D. Thus the assumption of constant density would be appropriate leading to simpler hydrodynamics, but the radiative transfer may have to be done as a time-dependent problem since t_B is less then the recombination or inonization time scale when $t_B \lesssim 1$ s. In other words the problem which needs to be solved for most flash phase EUV and Hα emission is the response of a static chromosphere to a sudden injection of energy by fast electrons balanced by radiative cooling.

While this has not been done, it has been shown by Kane and Donnelly (1971) that only 10^{28} erg are required for impulsive EUV emission in the 10–1030 Å range in a small flare and only 4.5×10^{25} erg is required for one Hα kernel (Vorpahl, 1973). Thus we may say that the dumping of only a small fraction of the accelerated electrons into a thick target in the low chromosphere is consistent with small flare energy requirements, but until the problem outlined above is solved, such checks are at best indicative.

5. Type III Radio Bursts and Interplanetary Electrons

A small fraction of the electrons accelerated which manage to move on open field lines give rise to type III bursts and interplanetary electrons. Although the energy involved may be small, the potential information content is large because these bursts together with their exciting electrons serve as electron probes of the corona from about 100000 km above the photosphere to the orbit of the Earth and an unknown distance beyond. The spectrum of interplanetary electrons as a function of time provides important clues on the acceleration of electrons during the flash phase.

The theory for type III bursts can be divided into three main parts: (1) The interaction of a spatially and temporally inhomogeneous electron distribution with a plasma and the resultant distribution of plasma waves. (2) The conversion of plasma waves into radiation. (3) The propagation of radiation from its source to the observer. For conciseness we refer to these as the plasma wave source, the radiation source and propagation. The theory has recently been reviewed in detail (Smith, 1974a) and a more complete account can be found there.

5.1. Plasma wave source

In the recent past this part of the problem has been the subject of great controversy largely because of incorrect or incomplete applications of theory. The observations of Stewart (1965) and Fainberg and Stone (1970) seemed to imply that the velocity of the exciting electrons was constant for several tens of solar radii. Application of the theory of quasilinear relaxation of an electron beam which is spatially homogeneous (Vedenov *et al.*, 1961; Drummond and Pines, 1962) showed that the beam should lose a significant fraction of its energy in at most 100000 km (Sturrock, 1964; Kaplan and Tsytovich, 1968; Melrose, 1970b; Smith, 1970a). In quasilinear relaxation the beam excites plasma waves with phase velocities slightly lower than the beam which in turn accelerate electrons with slightly lower phase velocities, etc. until a quasi-plateau is formed from the beam as shown in Figure 7. On the basis of this apparently large energy loss, various mechanisms were proposed to suppress this relaxation; i.e. to stabilize the beam, and the search is still in progress (Papadopolous *et al.*, 1974). The results of all attempts including that of Papadopolous *et al.* is that some rather special

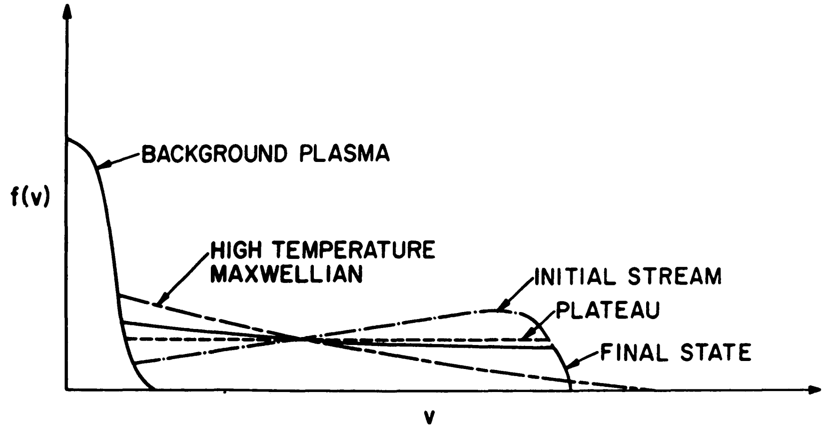

Fig. 7. Forms of the electron distribution function at various times in quasilinear relaxation. In a homogeneous plasma an initial stream (·–·–·) relaxes to a plateau (– – –) and finally to a high temperature Maxwellian (- - - - -) because of spontaneous emission processes. In the inhomogeneous relaxation characteristic of type III bursts there is insufficient time to reach a high temperature Maxwellian and the final state (———) is in between it and a plateau, closer to a plateau.

conditions have to be satisfied to stabilize an electron stream. For example, the nonlinear mechanism used by Papadopolous *et al.* is the oscillating two-stream instability which depends quite sensitively on the effective width of the stream since it only works well when the phases of all the waves are the same. Since we know from the previous sections that the stream must start out as a power law, it is likely to have a broad distribution at all times and hence fixed phase effects such as suggested by Papadopolous are unlikely to be important in the main part of the burst. Both Zheleznyakov and Zaitsev (1970a) and Smith and Fung (1971) have shown that the strongest nonlinear effect for a broad stream distribution, nonlinear Landau damping, is unable to stabilize an electron stream. Thus the stream must relax and the fact that the stream is both temporally and spatially inhomogeneous must prevent the rapid loss of energy indicated for a homogeneous stream.

Zheleznyakov and Zaitsev (1970a) were the first to consider this situation. They used a one-dimensional formalism and took collisional damping as well as Landau damping of the stream-plasma system in the direction of the stream into account. However, they integrated over the plasma wave spectrum in order to obtain a solution which is not allowable to correctly consider deceleration of the stream (Smith, 1973) or for the purpose of computing radiation near the fundamental of the plasma frequency (Smith, 1970a, 1974a). For these purposes the detailed shape of the plasma wave spectrum is important. In the low frequency regime (3 MHz and lower) Zaitsev *et al.* (1972) have constructed a one-dimensional similarity solution which only takes into account Landau damping of the stream in the direction of the stream. This solution also required integration over the plasma wave spectrum as well as complete neglect of the characteristics of the background plasma. This solution is completely lossless so that all of the energy given to the plasma waves by the stream is reabsorbed by the stream and the stream travels with constant velocity. This solution can be criticized on two major counts. (1) Although neglect of Landau damping by the background plasma may be allowable above 3 MHz because collisional damping is stronger, below 3 MHz Landau damping by the background plasma exceeds collisional damping (Harvey and Aubier, 1973) and must be included. Unfortunately a two-dimensional formalism for the plasma waves is required for this purpose (Smith, 1974a). (2) By integrating over the plasma wave spectrum the basic nature of quasi-linear relaxation which is to reduce the average velocity of the electrons involved has been neglected. As noted above, the manner in which the quasi-plateau in Figure 7 is reached is that fast electrons of velocity v_s excite plasma waves whose phase velocities v_{ph} are less than v_s. These plasma waves then accelerate some electrons to velocity $v_{\mathrm{ph}} < v_s$ which relax and produce plasma waves with phase velocities $v'_{\mathrm{ph}} < v_{\mathrm{ph}}$ and so on (Tsytovich, 1970). The resultant distribution of plasma waves (see, e.g. Nunez and Rand, 1969) thus has an average phase velocity v_{ph} considerably less than the initial average stream velocity v_s. Since the primary interaction of these waves is with electrons of the same phase velocity via the Cherenkov condition

$$\omega_p = \mathbf{k} \cdot \mathbf{v}, \tag{5.1}$$

where ω_p is the frequency of the plasma wave with wave vector $\mathbf{k}$ and $\mathbf{v}$ is the electron velocity, these waves cannot on the average give energy back to electrons with velocity v_s. In other words even in inhomogeneous quasilinear relaxation, the average velocity of the electrons must decrease. Since we now know that both the velocity of the leading edge and maximum of type III bursts do decrease (Fainberg *et al.*, 1972), this facet of quasilinear relaxation is consistent with the observations.

Because an investigation of the complete inhomogeneous quasilinear equations is likely to take some time, simple solutions in the spirit of Zaitsev *et al.* (1972), but which may provide a closer approximation to the real situation are desirable. For this reason Smith (1973) constructed a solution which overestimates the real losses suffered by a stream and which may be termed the 'completely lossy' solution. In this case which is an opposite limiting case to that of Zaitsev *et al.*, each time the stream forms a hump, the hump is cutoff and the valley filled in numerically as shown in Figure 8. It is assumed that the difference in energies of the initial and final distributions is released into plasma waves and none of the energy is returned to the stream. For an initial power law distribution of electrons with $\delta' = 2.3$ which is the value for the simultaneous electron-type III event analyzed by Lin *et al.* (1973), the deceleration of the leading edge of the burst from 10–20 $R_\odot$ can be compared with the average deceleration determined experimentally (Fainberg *et al.*, 1972) as shown in Figure 9. For comparison, the result of Zaitsev *et al.* (1972) would be a horizontal line in this figure. Because the scattering properties of the interplanetary medium for low energy electrons are unknown and some of the apparent deceleration could be due to scattering in pitch angle, whether the results of Smith (1973) satisfy the observations more closely than those of Zaitsev *et al.* (1972) is an open question. The main result to be drawn is that depletion of very fast electrons occurs quite rapidly and hence a more

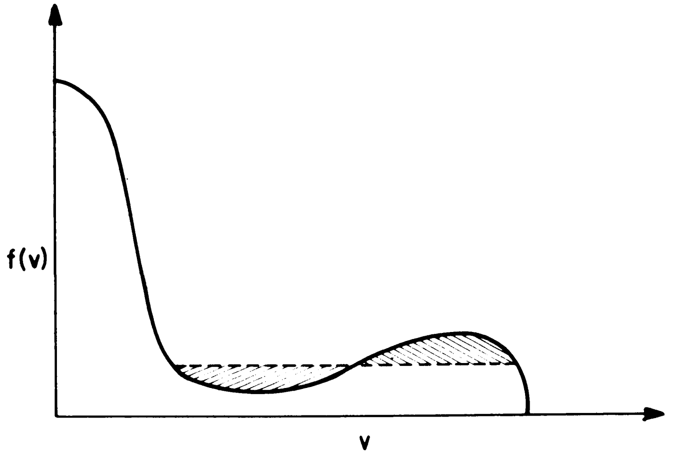

Fig. 8. Procedure of 'cutting off the hump and filling in the valley'. The original distribution (———) is converted into the final distribution (---) by cutting off the peak in the distribution (\\\) and filling in the valley (///) in such a way that the area cut off equals the area filled in.

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complete analysis taking into account losses more realistically as well as following the two-dimensional axially symmetric plasma wave distribution in detail is required. The stream was not followed beyond 20 $R_\odot$ for this reason. The interested reader may find the equations to be solved in Smith (1974a).

Beyond this basic stream-plasma interaction, many questions remain to be answer-

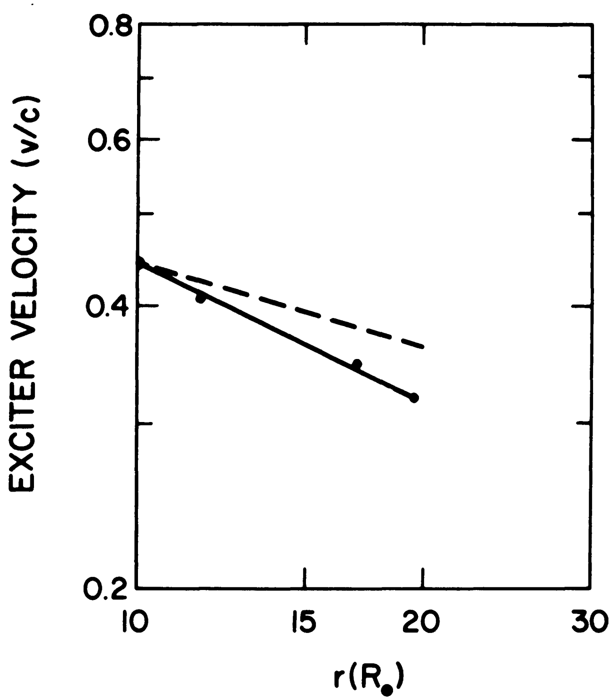

Fig. 9. Deceleration of the leading edge of a type III burst between 10 and 20 $R_\odot$ according to Smith's model (———) and according to the observations for a power law velocity distribution with index 4.6 and an initial stream density of 6.3×10^5 cm^{-3}, at a few thousand kilometers above the photosphere.

ed. What are the limitations of the above approach and when do nonlinear effects play a role? Although it takes about 100000 km for an initial power law to become unstable due to fast electrons overtaking slower ones (Smith, 1973), the starting frequencies of type III bursts are often much lower than the corresponding plasma frequency at this height (Stewart, 1974). A possible explanation for this is that quasi-linear relaxation may be suppressed by ion-acoustic waves or other inhomogeneities in the corona (Sturrock, 1964; Smith, 1970a, b). The occasional intermittency of type III bursts (Elgarøy, 1961; Ellis, 1969) argues that this must be possible. An alternate explanation for low starting frequencies is that electrons are accelerated high in the corona.

Another question is how does the stream become charge and current neutralized? This question was considered by Smith (1972a) for the case of a narrow velocity spread of the stream. While this is not allowable for the peak of the burst, it may be applicable to the leading edge of the burst. In the frame of the plasma, when the stream enters the plasma, the plasma reacts by accelerating electrons in the opposite direction to the stream to a velocity

$$v_a = \frac{-2n_s v_s}{n_e}, \tag{5.2}$$

where n_s is the density of the stream. As a result a current now flows in the direction opposite to the stream because the plasma has overreacted to the presence of the stream. The plasma then reacts to the presence of this current in the same manner as to the initial stream and accelerates electrons in the direction of the stream so that the initial stream is present which it reacts to again and so forth. These current oscillations which are accompanied by charge oscillations form a large amplitude plasma wave at the head of the stream which is damped by collisions. This large amplitude plasma wave may be related to the possible precursor phenomena observed by de la Noë and Boishot (1972) and others.

Finally there is the question of the macroscopic interaction of the electron stream with the inhomogeneous magnetic field and density structure of the corona. There is good evidence that electrons travel along coronal streamers in the metric band (Stewart, 1974), but no such evidence in the hectometer and kilometer band. Smith and Pneuman (1972) studied the motion of particles in a streamer model of Pneuman (1972) which included finite conductivity and added the curvature of the streamer due to solar rotation. In this model there is a small field across the streamer. They found that electrons could not travel more than about 1 $R_\odot$ in the current sheet along the streamer axis and have the properties required for type III bursts because of the finite transverse field. Thus the electrons causing type III bursts, which are most likely accelerated in current sheets as in Section 2, must drift out of these sheets due to the curvature of the streamer. Some aspects of the fine structure of type III bursts may also be due to the large scale inhomogeneous structure of the corona.

5.2. Radiation source

The basic processes for radiation near the fundamental and second harmonic of the plasma frequency are by now well known. Near the fundamental the process is the scattering of a plasma wave p on the polarization cloud of an ion i (electron density fluctuation) to produce a transverse electromagnetic wave t

$$p + i \rightarrow t + i', \tag{5.3}$$

with the polarization cloud going to a new state i' due to the absorption of some energy and momentum. Process (5.3) can be either spontaneous or induced and the absorption coefficient can be negative, i.e. the radiation can be amplified. It was shown by Smith (1970a) that the fundamental must be amplified to obtain observed ratios of fundamental/second harmonic power. Thus the fundamental is produced by a process which is effectively much more nonlinear than for the second harmonic. Zheleznyakov and Zaitsev (1970b) concluded that the fundamental would not be

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amplified using an inapplicable formula from Tsytovich (1970). The radiation which results from the amplification process has been examined in a detailed model of the source region at the 80 MHz level based on approximate plasma wave spectra (Smith, 1974b).

Near the second harmonic the radiation process is the combination of two excited plasma waves p and p' to produce a wave $t(2\omega_p)$

$$p+p'\rightarrow t(2\omega_p). \tag{5.4}$$

In contrast to process (5.3), the absorption coefficient for process (5.4) can only be positive. By considering this absorption, Melrose (1970a) and Smith and Sturrock (1971) showed that both of the combining plasma waves must be excited, i.e. have high effective temperatures, in order to obtain nonthermal radiation. Energy and momentum conservation require that the two plasma waves meet almost 'head-on'. Since the plasma waves produced by quasilinear relaxation are directed primarily in the direction of the stream, some secondary plasma waves must be produced moving in the opposite direction by scattering on the polarization clouds of ions

$$p+i\rightarrow p'+i'. \tag{5.5}$$

Again this process can be either spontaneous or induced. Generally the energy density in plasma waves is not sufficiently high for the induced process to play an important role. The emission and absorption coefficients for this case were calculated by Melrose (1970a) and Smith (1972b). The emission pattern for $v_s \lesssim 0.5\,c$ is a quadrupole.

At high frequencies (~ 100 MHz) the simultaneous observation of fundamental and second harmonic emission can provide important information about the source as shown by Smith (1970a). From the second harmonic power the energy density in plasma waves W_p required for a given source area A can be determined. By tracing rays in model sources of varying areas A with corresponding W_p and taking into account amplification of the fundamental, the area A and energy density W_p can be determined uniquely for a given source geometry. The results of such an analysis for a strong source at the 80 MHz level with $P(\omega_p)/P(2\omega_p)\approx 10$, where P is source emissivity, are that the source is about 14000 km or 0.3 min in diameter, that $W_p \approx 3\times 10^{-7}$ erg cm^{-3} and that all of the important rays come out the sides of the model cylinder travelling outward within an angular range of 26° (Smith, 1974b). However, scattering of the radiation in the source was neglected in deriving these results.

At low frequencies (~ 300 KHz) the energy density in plasma waves becomes so small that effective amplification of the fundamental is no longer possible and emission near the second harmonic becomes dominant. In this case there is no way to determine the source parameters from the radiation itself.

Melrose and Sy (1972) have recently considered the radiation mechanisms (5.3–4) in the presence of a magnetic field. For the case of a weak magnetic field, i.e. $\omega_{He}^2 \ll \omega_{pe}^2$, the only effect of the presence of the field is to give a slight or complete circular polarization to both fundamental and harmonic radiations. These results coupled with the observed degree of circular polarization could be used to set an upper limit to the magnetic field in the source if there were no mode coupling outside the source. However, the analysis of Dodge (1972) indicates that there may be significant mode coupling (see below).

5.3. Propagation

The radiation leaving the source is scattered, refracted, polarized and absorbed on its way to the observer. All of the recent work on this part of the problem has neglected the polarization of radiation and concentrated on the effect of scattering by random density inhomogeneities in a refracting absorbing medium. The method of treating this problem is to trace rays and to let them suffer random changes in direction after travelling a step size in a refracting absorbing plasma (Steinberg *et al.*, 1971; Riddle, 1972a, b). Since the results of this procedure have been summarized by Stewart (1974), and Smith (1974a) we only comment on some of Stewart's possible interpretations. Riddle (1974) has considered the case of anisotropic inhomogeneities in which the scale length $h = a\varrho$ in directions normal to the radius and $h = ab\varrho$ in the radial direction. With $b = 5$, Riddle found no significant differences for a point source compared with the isotropic inhomogeneity case. Riddle (1974) also showed that with scattering in a streamer taken into account the differences between isotropic emission, preferential emission backwards (Zheleznyakov and Zaitsev, 1970b) and preferential emission sideways (Smerd *et al.*, 1962) for the second harmonic would not be detectable observationally for a point source. Thus there is no reason to consider the possibility of preferential backward emission without considering a finite size source. It is quite possible that the intrinsic source size increases during a burst due to a larger number of field lines being 'lit-up' by high-energy electrons. It is also possible that the effective fundamental source is smaller than the second harmonic source because amplification by plasma waves is only effective in the center of the source where W_p is sufficiently high.

Consideration of finite size sources, time varying sources and different size fundamental and second harmonic sources appear as logical next steps for this part of the problem.

The observations of Dodge (1972) and Grognard and McLean (1973) on the polarization of type III bursts imply that the polarizing region must be a few solar radii above the radiation source at metric and decametric frequencies. This result is also in accord with that of Riddle (1974) who showed that the linear polarization observed is unlikely to represent conditions at the source because of the varied paths and path lengths by which radiation reaches the observer. The only known mechanism to convert the circularly polarized emission expected from the radiation source into primarily polarized radiation is mode coupling within regions of quasi-transverse magnetic field (Cohen, 1960; Zheleznyakov and Zlotnik, 1964). Cohen's mechanism for linear polarization requires that one of the incident magnetoionic modes be stronger than the other and the degree of linear polarization depends directly on the ratio of intensities. Thus strong linear or highly elliptical polarization would require that the quasi-transverse region be illuminated almost completely by o-mode radiation which is possible under the condition (Melrose and Sy, 1972)

$$v_{\mathrm{ph}} > \max\{v_e(3\omega_{pe}/\omega_{He})^{1/2},\, v_i(2\omega_{pe}/\omega_{He})\}, \tag{5.6}$$

where v_i is the ion thermal velocity.

5.4. Electron propagation in the interplanetary medium

Besides gradient drifts, etc. in large scale field inhomogeneities, electrons will also be scattered by magnetic field fluctuations. Unfortunately, analytic approximations are inapplicable in the energy range 10–100 keV so that no adequate theory for such propagation exists as for higher energy electrons (Jokopii, 1971). It appears that the best approach to the problem is a combination of numerical modelling and empirical results. Once the complete quasilinear equations (Smith, 1974a) which neglect pitch angle scattering are solved numerically, the amount of deceleration due to quasilinear relaxation will be known. Any additional apparent deceleration could then be attributed to scattering of electrons in pitch angle by interplanetary magnetic field fluctuations.

Kane and Lin (1972) and Lin (1973) have concluded that the electron acceleration region must have a density $n_e \lesssim 2\times10^9\ \mathrm{cm}^{-3}$ by noting that power law spectra observed at the earth have no turnover at low energies before ~6 keV. This deduction is based on the premise that the only interaction that electrons encounter in escaping from the Sun in a socalled 'scatter-free' event is the Coulomb interaction. However, we have seen in Section 2 and again in subsection 5.1 that other quasilinear and nonlinear interactions involving plasma waves are possible and that these can reconstitute a power law spectrum out of a non-power law spectrum. Moreover, for the nonlinear interactions of Section 2 very little radiation would be expected by the mechanisms of subsection 5.2 because the plasma waves are rapidly transformed to the region of k-space where the radiation produced has very low group velocities and so is absorbed before it has a chance to leave the source (Kaplan and Tsytovich, 1973). For the case where nonlinear interactions are not sufficiently strong it is possible that some of the decimetric emission (Kundu, 1965) is associated with such plasma turbulence. Thus there is no compelling reason for accepting the argument of Kane and Lin (1972) and Lin (1973), and aside from placing the electron acceleration region in the low corona where the magnetic field is still quite high for reasonable efficiency, its height is an open question. The numbers used in the first two sections of this review of $h \approx 10000$ km above the photosphere with $n_e \approx 3\times10^9\ \mathrm{cm}^{-3}$ and $B \approx 500$ G are simply guesses which satisfy the rather loose requirements stated above.

6. Discussion

We have reviewed our current knowledge of the mechanisms responsible for the nonthermal and impulsive quasi-thermal phenomena occurring during the flash phase of a flare. The relation of these mechanisms to the main thermal phenomenon going on during the flash phase, namely the rise of the soft X-ray burst is of interest. It is often observed that the rate of rise of the soft X-ray flux maximizes at the time of the most impulsive hard X-ray or microwave spike and that the flux continues to rise into the flare proper after the hard X-ray spike is over, but at a slower rate (Neupert, 1968). Thus the acceleration mechanism for electrons should be characterized by maximum heating being coincident with maximum acceleration. This is a natural consequence of the sequence of mechanisms proposed for acceleration in Section 2. Namely, the amount of energy going into heat in the reconnection process will be maximum when reconnection and hence acceleration is most efficient. The production of electron plasma and ion-acoustic waves by the accelerated particles will be most efficient when the acceleration is most efficient. The ion-acoustic waves in turn are heavily damped on a time scale $\tau_i \approx (m_i/4\pi n_i e^2)^{1/2}$ until T_e becomes much larger than T_i and the energy of the waves is rapidly converted into electron thermal energy. As noted in Section 2, the increase of electron temperature will tend to decrease the efficiency of acceleration by electron plasma waves. Since it will also make scattering through large angles more efficient, we know from the results of Friedman (1969) that the increase in electron temperature will make the direct first-order Fermi acceleration ineffective for all but a very select group of fast ions. In other words the process of rapid acceleration in a reconnection geometry is self-quenching by nature and this is consistent with the observations. However, both before and after efficient acceleration the current sheet is capable of acceleration of particles to energies slightly above their thermal values by slow reconnection. This energy is rapidly converted into heat since the collision mean free path of these slightly nonthermal particles is small, especially in a turbulent plasma. This fact allows an explanation for the fact that highly nonthermal phenomena such as type III bursts appear to rise out of a plasma of increasing temperature as indicated by soft X-rays (Teske *et al.*, 1971) and that the soft X-ray flux continues to rise after the impulsive phenomena are over (Neupert, 1968).

If we now ask what needs to be done to improve the theories of flash phase mechanisms, it is not hard to find several good problems. The parts of the acceleration mechanism outlined in Section 2 need to be worked out in detail. To start with, what determines the rate of reconnection? Can the acceleration of particles act like a siphon to suck field lines and more particles into the acceleration region or is everything determined by conditions far from the reconnection region? What is the spatial distribution of plasma turbulence as a function of time and can it provide power law spectra with the indices observed? What are the fates of particles brought in on incoming field lines? How does the temporal development of plasma turbulence affect the efficiency of the first-order Fermi process?

In the hard X-ray interpretations are we being too naive in favoring dismissing the possibility of a significant thick-target contribution in most flares? Better observations of the location of the hard X-ray source will help to answer this question. In the interpretation of impulsive microwave bursts can we find some relatively simple means of determining which absorption mechanism is dominant in a given situation? The impulsive EUV and Hα emission during the flash phase is perhaps the problem most ripe for solution. A large amount of data has recently become available (Kane, 1974), but there is no consistent model to which it can be fit. The construction of such a model will inevitably involve a detailed calculation of the radiative response of the chromosphere in hydrogen to a flux of impulsive nonthermal electrons as outlined in Section 4. The problem of impulsive EUV lines and continuum of atoms other than hydrogen can then be attacked since the hydrogen lines and continuum determine the basic state of the chromosphere.

While the type III burst problem has probably seen more progress than other problems because of its relative simplicity, much remains to be done. The most pressing problem is the numerical solution of realistic quasilinear equations which are inhomogeneous in time and space. With this problem solved and the resultant realistic plasma wave spectra, the radiation source can be defined much better. With a better radiation source, the problems of scattering and polarization of radiation can be treated more realistically, etc. With a complete numerical solution of the quasilinear equations, the amount of apparent deceleration due to pitchangle scattering and thus the diffusion coefficients for low-energy electrons in the interplanetary medium can be determined. These can then be used to treat the stream-relaxation more realistically and thus refined by successive approximations. With this information in hand, we could define the required electron source much more precisely.

For conciseness and to avoid too much speculation this review has followed a rather narrow line. Certainly many other phenomena are most likely initiated at the flash phase such as the white light and CN flare, the formation of type II burst sources and interplanetary shock waves, the acceleration of particles to relativistic energies and possibly the formation of type IV burst sources. There are also many known interrelations between various phenomena such as a relation between Hα absorption features and type III bursts (Axisa *et al.*, 1973; Kuiper and Pasachoff, 1973). However, only rudimentary qualitative ideas exist for many of these phenomena and the rest are more properly treated along with the main phase of flares.

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DISCUSSION

Wild: Can you account for the quasi-periodicity of groups of type III bursts?

Smith: No, I can't in this simple model. However, one can imagine a sheet with some extent and a hydromagnetic wave moving down the sheet causing reconnection to be enhanced as it passes, which might explain the quasi-periodicity.

Mayfield: Can you give an estimate of the lifetime of these current sheets?

Smith: No, because I haven't studied their formation and we don't know what determines the rate of reconnection which will lead to their decay.

Lin: You used 500 G in your estimates. This seems like a high value for a height of 10000 km.

Smith: As I noted the acceleration is proportional to B^2. I used 500 G because if we have a dipole 10000 km long of strength 500 G, then it will still be 500 G at a height of 10000 km.

Dryer: Could a jump in the magnetic field (as in a fast MHD shock) produce accelerated electrons which could produce type III bursts?

Smith: No, the field should be antiparallel as in a separate body of plasma with its own field interacting with the ambient field.

Dryer: As in a plasma piston, perhaps?

Smith: Yes.

Rosenberg: Kuperus and I looked into neutral sheet configurations. Taking a flux tube and twisting its end points, it becomes kink unstable, but stabilizes nonlinearly with a kink. Going on, you end up with a narrow, very long spiralled neutral band (giving the possibility for fast reconnection). The other advantage is the great amount of magnetic energy present with relatively small field strengths.

Smith: I have treated the simplest configuration here. One of the problems with this configuration is that you need a large area of sheet to release a significant amount of energy. Any way you can twist up a sheet into a three-dimensional configuration is likely to be better than the configuration in my talk.

Brown: I have two comments: (1) I agree with your criticisms of the radiative loss model in my 1973 paper but as regards the question of constant pressure or density, I considered the former because it is appropriate in large events, e.g. Cline *et al.* (1967) which *e*-folded in 100 s as against a 50 s dynamical time. When you modify my analysis to constant density, the higher density greatly enhances the radiative loss (even neglecting Hα) and it becomes hard to get the flare hot enough. (2) It is important to point out that models of flare heating by electrons, widely mentioned here today, strongly depend on how low the steep electron spectra extend in energy since all their energy resides at the low end. This limit is not merely instrumentally hard to determine, but is mixed with the thermal X-ray contribution which must therefore be considered carefully, as I hope to do in Part III.

Smith: I agree with you completely on your second comment. Regarding your first comment, it still seems to me that whether constant density or constant pressure is appropriate depends on how much energy is delivered how quickly. It is quite possible that in small flares we are seeing essentially direct conversion of electron energy into EUV radiation with no significant heating of the chromosphere.

Kane: In Takakura's model the energy is transported from the acceleration region to the EUV and optical source through heat conduction. What time constant do you expect for such an energy transport?

Smith: I believe the fine structure in impulsive EUV and hard X-ray data rule out any heat conduction model because the time constant would be too long.

McLean: With regard to your explanation of linear polarization in type III bursts, I should like to ask if you do not find this inconsistent with the observation of circular polarization in type I emission?

Smith: No, because the two types of bursts are thought to be produced in different regions of the corona. Melrose might wish to comment.

McLean: The problem is made more difficult by the fact that these bursts are seen over a wide range of longitude and so we should certainly expect to see both types of burst through the same parts of the corona, yet for one type of burst circular polarization is converted into linear and not for the other.

Melrose: It is thought that the difference between type I and type III is associated with the much stronger field strengths for the former. If one wishes to explain linear polarization in type III, but none in type I, by mode coupling at a *QT* region, this region must be low in the corona because otherwise *both* type I and type III would be linearly polarized. However with a *QT* region low in the corona the Faraday rotation between the *QT* region and Earth should wash out the linear polarization. Linear polarization by this mechanism seems unacceptable.

Stewart: How do you explain the high level in the corona at which type III's start in type I–type III storms?

Smith: There are two possibilities: (1) The electrons are generated low, but because the corona is turbulent to a certain level, generation of plasma waves is supressed. (2) Electrons are accelerated in current sheets high in the corona.

Krall: There is a recent result from Papadopoulos *et al.* which provides an alternate mechanism from either Zaitsev's or Tsytovich's for stabilizing an electron beam. This mechanism is based on the fact that electron plasma waves in the presence of a beam can decay unstably into waves which are not resonant with the beam. This prevents the formation of the quasilinear plateau (Zaitsev's idea) as well as giving an ion-wave coupling stronger than that calculated in Tsytovich's model.

Smith: I am familiar with the work of Papadopoulos *et al.* on stabilizing a narrow stream by the oscillating two-stream instability. However, this is a fixed-phase effect which wouldn't work for a wide stream as expected from a stream which starts out from a power law. I think using the simplest approach is best unless it can be shown to lead to inconsistent results, i.e. quasilinear relaxation.