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Particle Acceleration Processes in the Solar Corona

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Abstract

Theoretical ideas on particle acceleration associated with solar flares are reviewed. A historical outline is used to introduce the various acceleration mechanisms. These are stochastic acceleration in its various forms, diffusive acceleration at shock fronts, shock drift acceleration, resonant acceleration, acceleration during magnetic reconnection and acceleration by parallel electric fields in double layers or electrostatic shocks. Particular emphasis is placed on so-called first phase acceleration of electrons in solar flares, which is conventionally attributed to bulk energisation of electrons (Ramaty *et al.* 1980). There is no widely accepted theory for bulk energisation, which may be regarded as an enhanced form of heating. Ideas on bulk energisation are discussed critically. It is argued that the dissipation cannot be due to classical resistivity and involves anomalous resistivity or hyperresistivity, e.g., in multiple double layers. The dissipation must occur in very many localised regions. Bulk energisation due to magnetic reconnection is discussed briefly. A model for bulk energisation due to the continual formation and decay of weak double layers is outlined.

Other aspects of particle acceleration associated with solar flares are reviewed more briefly. The specific topics discussed include the following:–

(1) prompt acceleration of ions to $\gtrsim 30\,\mathrm{MeV}$ per nucleon on a time scale $\lesssim 1\,\mathrm{s}$, implied by data on prompt γ-ray bursts,
(2) second phase acceleration, emphasising radio evidence for at least six phases of electron acceleration as well as the data on solar cosmic rays,
(3) the problem of creating a seed population of energetic particles, required for the favoured second phase acceleration mechanisms to operate, and
(4) the preferential acceleration of ions of different species in connection with the so-called anomalous abundances in solar cosmic rays.

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Contents

1. Introduction

Acceleration of energetic particles in association with solar flares was originally separated into two phases (Wild, Smerd and Weiss 1963, de Jager 1969). Radio data suggest that there are many more than two phases of acceleration (Smerd 1975, Melrose and Dulk 1987), and there is unambiguous evidence for non-flare-associated acceleration, most notably in storms. However, it remains helpful to concentrate on two general classes of acceleration similar to the two phases as originally proposed. One class of accleration is bulk energisation in which, it appears, all the electrons in a substantial volume have their energies increased by a substantial fraction in a short time. The other class of acceleration consists of all those processes that involve suprathermal particles with a clearly nonthermal distribution. These include all the traditional second phase acceleration processes, such as acceleration by hydromagnetic turbulence and by shock waves.

The concept of bulk energisation of electrons arose in connection with the acceleration of electrons in the impulsive phase of solar flares. A working definition of *bulk energisation* in the impulsive phase is that it involves an increase in the mean energy of the electrons by a factor $\gtrsim 10$ in a time $\lesssim 1$ s. In bulk energisation, the energised electrons have a quasi-thermal distribution in the statistical sense that the ratio of the variance to the mean square energy is close to that for a Maxwellian distribution. The production of a nonthermal tail on the electron distribution and any associated bulk energisation of the ions are important questions that should be addressed in association with any detailed discussion of bulk energisation of electrons.

The hard X-ray data (e.g., Dennis 1985) are interpreted in terms of precipitating electrons emitting bremsstrahlung, and the data then require that $\gtrsim 20$ keV electrons precipitate at a rate up to about $10^{36}\,\mathrm{s}^{-1}$ (Hoyng, Brown and van Beek 1976). The energy involved in these electrons is a substantial fraction of the energy released in a flare. The most widely favoured interpretation is that the energy released goes primarily into such electrons for most flares, with a substantial energy going into mass motions and into a blast wave only for large flares. The rate ($\approx 10^{36}\,\mathrm{s}^{-1}$) of precipitation and the total number of precipitating electrons required (up to $\approx 10^{39}$) are so large that they lead to severe constraints on the model. These and associated constraints have led some authors to explore the alternative hypothesis that the hard X-rays are due to precipitating ions (Colgate 1978, Simnett 1986). Current opinion favours a dominant role for the electrons. It is relevant to note that many type III events observed in the interplanetary medium are associated with storms and not with flares (e.g., Suzuki and Dulk 1985, Kai, Melrose and Suzuki 1985); moreover, many such events are dominated by electrons of quite low energy (2–10 keV) that could not possibly come from a flare site due to the prohibitive collisional losses in propagating through the lower corona (e.g., Lin 1985). It follows that bulk energisation is not restricted to flares, and that it also occurs in storms where it is required to account for the electrons in storm type III bursts, and also presumably for the electrons that produce type I bursts. There is no consensus on the detailed mechanisms involved in bulk energisation.

There is a wide variety of acceleration mechanisms for suprathermal particles. These include stochastic acceleration of various kinds, including acceleration in neutral sheets, and acceleration by shock waves. (Stochastic acceleration is often called Fermi acceleration, but this can be ambiguous because 'Fermi' acceleration is used with two different meanings, one being the specific mechanism proposed originally by Fermi, cf. §2, and the other being the generic sense in which it is effectively synonymous with 'stochastic' acceleration.) A characteristic feature of most stochastic acceleration mechanisms is that there is a threshold energy below which any specific mechanism is ineffective or inoperative. Thus such acceleration is effective only when a seed population of energetic particles already exists. Put in other words, an injection spectrum of suprathermal particles is required. The acceleration increases the *energy* of the already suprathermal particles but does not increases the *number* of suprathermal particles. The production of suprathermal particles, that is, the formation of a nonthermal tail on a thermal distribution of particles, is regarded as a separate but essential problem.

The main emphasis in this review is on acceleration associated with solar flares. There is an extensive literature on this topic; some more recent reviews include those by Ramaty *et al.* (1980), Heyvaerts (1981), Forman, Ramaty and Zweibel (1985), de Jager (1986), Somov (1986*a*), Vlahos *et al.* (1986), Ramaty and Forman (1987) and Sakai and Ohsawa (1987). Related reviews are those of acceleration in the interplanetary medium, e.g., Fisk (1979), Pesses, Decker and Armstrong (1982), Forman and Webb (1985), and also some more general reviews of acceleration with some emphasis on solar applications, e.g., Toptygin (1980), Axford (1981), Toptyghin (1983).

In §2 acceleration mechanisms are introduced in a primarily historical context and then classified. Some existing and new ideas on bulk energisation of electrons are discussed in §3, and the associated (in the impulsive phase) prompt acceleration of ions is discussed in §4. Some general remarks on second phase acceleration mechanisms are presented in §5 and the important associated question of the production of suprathermal particles, needed as the seed population, is discussed in connection with so-called anomalous abundances in §6.

2. Historical Review

It is convenient to separate acceleration mechanisms into five general types: stochastic, shock, resonant, reconnection and parallel electric. The ideas behind each of these are introduced here from a historical viewpoint.

Stochastic Acceleration

The development of current ideas on the acceleration of particles is generally regarded as starting with the suggestion by Fermi (1949, 1954) that galactic cosmic rays are accelerated by bouncing off magnetised clouds. The specific mechanism proposed by Fermi is ineffective in practice, but the proposed mechanism contains ideas that are relevant to all versions of stochastic acceleration. One can identify three important ingredients in Fermi's mechanism. One is that a particle gains energy in a head-on collision with a cloud and loses energy in an overtaking collision. The second idea is that if the cosmic rays are moving at random then they have a higher probability of having a head-on rather than an overtaking collision. The first order changes in the energy in head-on and in overtaking collisions cancel in a statistical treatment, but the second order effects do not and the difference leads to a net average acceleration over a time long compared with the collision time. The third ingredient is more subtle because it is implicit: the *assumption* that the distribution of particles remains isotropic requires a specific mechanism that isotropises the particles. The reason is that the collisions tend to align the particle velocities along the magnetic field lines, and this alignment severely limits the acceleration.

Fermi-type acceleration also occurs in the so-called betatron effect (magnetic pumping in the plasma physics literature) in a magnetic trap when there are temporal variations in the strength B of the magnetic field, in which case conservation of the magnetic moment implies that the energy of the particle is proportional to B (Swann 1933, Schlüter 1957, Berger *et al.* 1958). A related process is transit acceleration (e.g., Shen 1965) when particles diffuse through an inhomogeneous B. The important features of all such mechanisms is that they involve statistical energy gains and that they require a mechanism that tends to maintain the assumed isotropy of the particles.

These ideas are applicable to acceleration by MHD turbulence. Qualitatively, the efficiency of the acceleration is higher when the energy changes are both frequent and small than when they are both infrequent and relatively large, as in Fermi's mechanism. In earlier treatments of this effect (Thompson 1955, Kaplan 1956, Davis 1956, Parker 1957, Parker and Tidman 1958), the need to maintain isotropy was implicit. Various mechanisms to maintain isotropy were included in later statistical treatments (e.g., Asséo and Barthomieu 1966, Sturrock 1966, Hall and Sturrock 1967), with the main emphasis being on resonant scattering (e.g., Hasselmann and Wibberenz 1968, Melrose 1968*a*, Kulsrud and Ferrari 1971). The importance of resonant scattering was first recognised in connection with particles trapped in the earth's radiation belts; the scattering of ions involves Alfvén waves (Wentzel 1961, Dragt 1961) and the scattering of electrons involves whistlers (Dungey 1963, Cornwall 1964), cf. also Kennel and Petschek (1966). These waves can be generated by the particles themselves due to their induced anisotropy (Melrose 1974). In this mechanism, the low frequency MHD turbulence provides the energy and the high frequency waves maintain the isotropy.

The nature of stochastic acceleration was further clarified by Achterberg (1981) who pointed out that the interaction with the low frequency turbulence may be interpreted in terms of a resonance at harmonic number $s = 0$, whereas the scattering by the high frequency waves is due predominantly to resonances harmonic numbers $s = \pm 1$. An important implication of Achterberg's approach is that Fermi-type acceleration is due entirely to magnetoacoustic turbulence. The point is that MHD turbulence in general consists of a mixture of two modes, both of which propagate at approximately the Alfvén speed v_A in a strongly magnetised plasma. (There is a third mode, the slow mode, which is not relevant to the discussion here.) The magnetoacoustic or fast mode is compressive and propagates almost isotropically, and the Alfvén mode is tortional and has an energy flow only parallel to the field lines. One class of model for acceleration by MHD turbulence involves a magnetic cascade, e.g., Bicknell and Melrose (1982), van Ballegooijen (1986). In a magnetic cascade there is a source of MHD turbulence with a large characteristic scale length. The resulting long-wavelength turbulence breaks up into shorter wavelength turbulence, with a dissipation-free flow of energy in $\boldsymbol{k}$-space from the source at small $|\boldsymbol{k}|$ (long wavelengths) through the so-called inertial regime to a sink at large $|\boldsymbol{k}|$ (short wavelengths). Each break-up should produce a roughly equal mixture of the two modes irrespective of the initial mode of the MHD turbulence. However, only the magnetoacoustic component damps due to Fermi-type acceleration of fast particles. Alfvén turbulence is ineffective in accelerating particles through a Fermi-type process because the matrix element for resonance at $s = 0$ is very small for Alfvén waves in comparison to that for magnetoacoustic waves at the same wavelength.

The reason the energy of particles should increase in general in stochastic acceleration can be understood from a thermodynamic-type argument. In Fermi's mechanism one may define an effective temperature for the turbulence by equating the effective thermal energy to the mean kinetic energy of the clouds. This temperature is enormous, as is the effective temperature for MHD turbulence when it is defined. Then given a mechanism that allows exchange of energy between the turbulence and individual particles, the effective temperature of the particles tends towards equalisation with that of the turbulence. Thus the mean energy of particles tends to increase towards the very large effective thermal energy of the turbulence.

A characteristic feature of stochastic acceleration is that the average rate of acceleration may be described by a diffusive equation in momentum space. This is of the form (Tverskoi 1967, 1968, Kulsrud and Ferrari 1971)

$$\frac{\partial}{\partial t}\langle f\rangle(p) = \frac{1}{p^2}\frac{\partial}{\partial p}\left[p^2 D_{pp}(p)\frac{\partial}{\partial p}\right]\langle f\rangle(p), \tag{2.1}$$

where $\langle f\rangle(p)$ is the particle distribution function averaged over pitch angle. For Fermi-type acceleration, the diffusion coefficient is of the form

$$D_{pp}(p) = \zeta_A \frac{cp^2}{4v}\left(1 - \frac{v_A^2}{v^2}\right)^2, \tag{2.2}$$

where ζ_A is the acceleration rate. The final factor in (2.2), which was derived by Kulsrud and Ferrari (1971) and Achterberg (1981), does not appear in simpler treatments of Fermi-type acceleration; this factor may be ignored only in the limit of particle speeds much greater than the Alfvén speed, i.e., for $v \gg v_A$. The meaning of ζ_A is most easily understood by estimating the mean rate of acceleration (Tsytovich 1966, Melrose 1968*a*)

$$\left\langle\frac{d\varepsilon}{dt}\right\rangle = \frac{1}{p^2}\frac{\partial}{\partial p}\left[vp^2 D_{pp}(p)\right]. \tag{2.3}$$

In the limit $v \gg v_A$ (2.3) gives

$$\left\langle\frac{d\varepsilon}{dt}\right\rangle \approx \zeta_A pc, \tag{2.4}$$

which reduces for highly relativistic particles to the familiar form $\langle d\varepsilon/dt\rangle \approx \zeta_A\varepsilon$. The acceleration rate may be estimated as

$$\zeta_A = \frac{\pi}{4}\bar{\omega}\left(\frac{\delta B}{B}\right)^2, \tag{2.5}$$

where δB is the magnetic amplitude in the waves and $\bar{\omega}$ is their mean frequency. A remarkable feature of (2.5) is that the acceleration rate is independent of the details of the scattering of the particles, and yet effective scattering is an essential ingredient in the theory. The parameters that appear in (2.5) refer only to the MHD turbulence that is causing the acceleration.

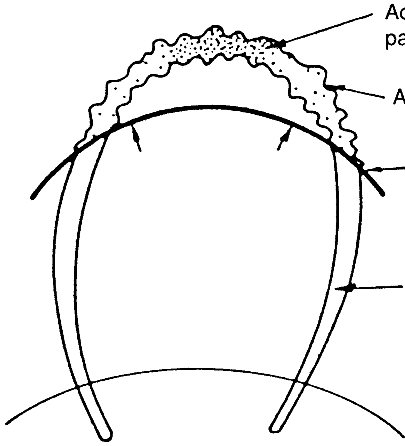

Fig. 1. The model proposed by McLean *et al.* (1971), as drawn by Wild and Smerd (1971), for a particular form of first order Fermi acceleration where the accelerated particles are trapped ahead of an advancing shock front and gain energy at every reflection from the front.

Acceleration at Shock Fronts

A different modification of the Fermi mechanism involves application to acceleration at shock fronts, as suggested by Parker (1958), Hoyle (1960), Schatzman (1963), Wentzel (1963, 1964) and Sonnerup (1969), and by many subsequent authors, e.g., Achterberg and Norman (1980) and others cited below in more specific contexts. An important idea is that of *first order Fermi acceleration* when particles reflect from two converging shock fronts. A specific model for this, cf. Fig. 1, seems particularly favourable for the acceleration associated with shock waves following solar flares (McLean *et al.* 1971, Wild and Smerd 1971). There is now an extensive literature on acceleration at shock fronts, especially in connection with (i) the earth's bow shock and interplanetary shocks and (ii) the acceleration of galactic cosmic rays and of relativistic electrons in synchrotron sources. More recent general reviews include those by Toptygin (1980), Axford (1981), Forman and Webb (1985), and reviews more specifically directed to solar flares include those by de Jager (1986), Ramaty and Forman (1987). It is now recognised that there are two quite different mechanisms by which a shock wave can accelerate particles. These are called *shock drift acceleration* and *diffusive shock acceleration.*

Shock drift acceleration occurs in a single reflection or transmission at a shock front, e.g., Toptygin (1980), Pesses (1981). This energy change is attributed to the electric field in the shock front, and its magnitude is determined by the potential energy associated with the electric field. The magnitude of the change is determined by the strength of the shock. As a consequence, the change in the energy of a particle is roughly independent of the initial energy of the particle. Such changes are important only when the initial energy of the particle is comparable with the change in energy. Thus shock drift acceleration is ineffective for highly energetic particles.

The details of shock drift acceleration follow relatively simply from an approximate conservation of the familiar first adiabatic invariant when a particle crosses the shock. Let the upstream and downstream regions be labelled 1 and 2, respectively. Using

$$\left(\frac{p_\perp^2}{B}\right)_1 = \left(\frac{p_\perp^2}{B}\right)_2, \tag{2.6}$$

the following give the maximum changes in the energy for reflected r and transmitted t particles compared with the incident i particles in the nonrelativistic limit (Toptygin 1980)

$$\text{Reflection } (1 \to 1): \qquad \frac{\epsilon_r}{\epsilon_i} = 1 + 4\left(\frac{B_2}{B_1} - 1\right), \tag{2.7}$$

$$\text{Transmission } (1 \to 2): \qquad \frac{\epsilon_t}{\epsilon_i} = 1 + 2\left(\frac{B_2}{B_1} - 1\right), \tag{2.8}$$

$$\text{Transmission } (2 \to 1): \qquad \frac{\epsilon_t}{\epsilon_i} = 1 + \left(\frac{B_1}{B_2}\right)^{1/2} + \left(1 - \frac{B_1}{B_2}\right)^{1/2}. \tag{2.9}$$

The maximum value of B_2/B_1 across a shock is 4.

The changes on reflection and transmission also alter the pitch angle of the particle. For example, in the model of converging shocks illustrated in Fig. 1, in the absence of scattering the particle velocities become increasingly aligned along the magnetic field as their energy increases, and this reduces the probability of their being reflected rather than transmitted on encountering a shock front. Thus the acceleration tends to be self limiting in the absence of scattering. In the presence of resonant scattering first order Fermi acceleration between converging shocks is a very effective acceleration mechanism.

Current ideas on diffusive shock acceleration may be regarded as a modification of this first order Fermi mechanism involving converging shocks. It was recognised in several different contexts that the high efficiency of acceleration also applies to a single shock provided that there are scattering centres both upstream and downstream of the shock (Krimsky 1977, Axford, Leer and Skadron 1977, Bell 1978*a,b*, Blandford and Ostriker 1978). The idea is that when viewed from a frame at rest in the fluid on one side of the shock, the scattering centres on the other side of the shock are moving towards the shock front. Hence, each time a particle crosses the shock front and is scattered, the first scattering is head-on and causes a net gain in energy. Further scatterings cause the particles on that side of the shock to maintain an isotropic distribution in the rest frame of the scattering centres there. A particle that crosses the shock in the opposite direction similarly has its first scattering head-on, and so the acceleration can proceed through many shock crossings for a typical particle. A major success of this theory is that the predicted spectrum of accelerated particles has a power law form of just the type observed in cosmic rays and synchrotron sources.

Resonant Acceleration

A resonant interaction between a particle with velocity $\boldsymbol{v}$ and a wave with frequency ω and wave vector $\boldsymbol{k}$ in an unmagnetised plasma occurs when the resonance Doppler condition

$$\omega - \boldsymbol{k} \cdot \boldsymbol{v} = 0 \tag{2.10}$$

is satisfied. The condition (2.10) corresponds to the wave having zero frequency in the rest frame of the particle. The electric field of the wave can accelerate the particle freely until the velocity has changed sufficiently so that the resonance condition (2.10) is no longer satisfied. The resulting efficient transfer of energy between waves and particles is from the waves to the particles for particles slightly slower than implied by (2.10) and from the particles to the waves for particles slightly faster than implied by (2.10). That is, the sign of the energy transfer is determined by the sign of $dF(v)/dv$, where $F(v)$ is the one-dimensional distribution function along the direction $\boldsymbol{k}$ of propagation of the waves. This resonant interaction underlies (i) collisionless or Landau damping (Landau 1946) of waves in a thermal plasma, (ii) the generation of Langmuir waves in a streaming instability or other situation where there are more faster particles than slower particles ($dF(v)/dv > 0$), and (iii) *resonant acceleration* of suprathermal particles by waves with phase speed $v_\phi = \omega/k$ equal to v when $\boldsymbol{k}$ and $\boldsymbol{v}$ are parallel.

The resonance condition in the presence of a magnetic field is

$$\omega - s\Omega - k_\parallel v_\parallel = 0, \tag{2.11}$$

where $k_\parallel$ and $v_\parallel$ denote components along the magnetic field, where $\Omega = |q|B/m\gamma$ is the gyrofrequency of a particle with charge q, mass m and Lorentz factor $\gamma = [1 - v_\perp^2/c^2 - v_\parallel^2/c^2]^{-1/2}$, and where $s = 0, \pm 1, \pm 2, \ldots$ is the harmonic number. In this case the sense of the energy transfer is from the waves to the particles for

$$\left[\frac{s\Omega}{v_\perp}\frac{\partial}{\partial p_\perp}+k_\parallel\frac{\partial}{\partial p_\parallel}\right]f(p_\perp,p_\parallel)<0, \tag{2.12}$$

and from the particles to the waves when inequality (2.12) is reversed.

The idea that such resonant interactions might be important in accelerating particles in astrophysical plasmas was suggested by Tsytovich (1963, 1965, 1966), cf. also Melrose (1968*a*). An obvious problem with resonant acceleration in specific applications is the source of energy for the waves. Resonant acceleration is an alternative to stochastic acceleration, in which case the energy is supplied in the form of MHD turbulence, e.g., Barbosa (1979), Eilek and Henricksen (1984). The formal distinction between them is that the resonant acceleration involves high harmonics (e.g., Lacombe 1977) and stochastic acceleration involves $s = 0$ in (2.11) (Achterberg 1981). A process that is relevant to this type of resonant acceleration is the *cascade of magnetic energy*, mentioned above in connection with Fermi-type acceleration. The cascade transfers turbulent energy from long wavelengths, where it can be generated effectively by mass motions, to short wavelengths, where it can be dissipated by accelerating fast particles or heating the plasma, e.g., Bicknell and Melrose (1982), van Ballegooijen (1986). In such a cascade effectively all the energy injected at the largest scale length ends up in the energetic particles provided that the dominant dissipation mechanism (the sink at short wavelengths) is determined by a dissipation process that involves transfer of energy to fast particles. This dominant mechanism then determines the dissipation scale length at which the cascade terminates and no other dissipation mechanism that might be important at shorter wavelengths is relevant. As remarked above, only the magnetoacoustic component is damped by Fermi-type acceleration, but in a cascade one expects that even if at some stage there is pure Alfvénic turbulence then at the next stage in the cascade it will produce a roughly equal mixture of the two modes at shorter scale lengths; it then follows that, as the magnetoacoustic component dissipates rapidly, ultimately all the turbulent energy is transferred to the energetic particles. Both modes can damp through resonant acceleration. Which of Fermi-type or resonant acceleration (or any other dissipation process) is the more important is determined by the one which would be effective at the longest wavelength. That is, if one determines the wavelength at which each dissipation process would be an effective sink terminating the cascade, then the dominant one is that which corresponds to a sink at the longer wavelength.

The case of most interest to solar flares is the acceleration of electrons by Langmuir waves. It is reasonable to neglect the ambient magnetic field in treating the wave-particle resonance and to use the resonance condition (2.5) rather than (2.11). One of the first detailed models for such acceleration was that of Lacombe and Mangeney (1969) which they applied type IV bursts. The source of the energy in the Langmuir case is the main difficulty with this mechanism. One idea that caused considerable interest was a particular mechansism for upscattering of ion sound waves into Langmuir waves, as first proposed by Tsytovich, Stenflo and Wilhelmsson (1975) and called *turbulent* *bremsstrahlung*. This process was of considerable interest because ion sound waves are readily generated in a plasma due to a current instability, and their upconversion would provide a source of Langmuir turbulence under conditions that plausibly obtain whenever magnetic energy is dissipated, e.g., Hoyng (1977*a,b*). The nature of this upconversion mechanism caused some controversy due to confusion with other possible upconversion mechanisms, cf. the literature cited by Kuijpers and Melrose (1985). Although both Kuijpers and I had independently discussed turbulent bremsstrahlung without questioning it (e.g., Kuijpers 1980, Melrose 1982), we have subsequently argued that the mechanism does not exist (Melrose and Kuijpers 1984, Kuijpers and Melrose 1985), and that its supposed existence resulted from an algebraic-type error in its original derivation (Melrose and Kuijpers 1987). Thus this particular upconversion mechanism should not be considered further. Although other upconversion mechanisms are possible, none seems at all favourable, and generation of the Langmuir turbulence from ion sound turbulence does not appear to be a viable source of energy for acceleration of the fast electrons associated with solar flares and storms.

Acceleration during Magnetic Reconnection

It is widely accepted that the energy released in solar flares is the result of magnetic reconnection or annihilation. An early theory for this was called 'discharge theory' by Giovanelli (1946, 1947, 1948). The discharge theory was criticised by Cowling (1953), who pointed out, *inter alia*, the importance of the formation of current sheets. Further development of electrodynamic models for solar flares, involving what is now called current dissipation or magnetic annihilation, followed in the work of Sweet (1958), Dungey (1958) and Gold and Hoyle (1960).

These early ideas were reviewed by Sweet (1969), cf. also Parker (1963), who emphasised the difficulties implied by the supposed slowness of magnetic reconnection due to the high electrical conductivity σ or low resistivity $\eta = 1/\sigma$ of the plasma. A simple model of a site of magnetic reconnection has the magnetic field along the z axis, and its amplitude $B(x)$ passes through zero as a function of x. This zero defines a neutral plane. The current implied by curl $\boldsymbol{B}$ tends to collapse to a current sheet in the neutral plane. Dissipation is due to the finite electrical resistivity of the plasma in the neutral plane which allows the magnetic field to diffuse through the plasma. The rate of this diffusion is slow due to the high electrical conductivity, and this seemingly limits the rate of energy release to an unacceptably slow rate for a flare model. The inclusion of anomalous resistivity increases the thickness of the current sheet and hence the rate of reconnection (e.g., Coppi 1983). A related problem is the removal of the incoming plasma from the sheet after annihilation has occurred. The latter problem was alleviated by the model for magnetic reconnection proposed by Petscheck (1964) who showed that slow shock waves can form to remove the processed plasma, cf. Fig. 2. Reconnection models as applied to both solar and magnetospheric physics were reviewed by Vasyliunas (1975).

The rate of magnetic reconnection in the corona cannot be observed directly. One possible way of estimating a plausible rate is to argue by analogy with estimates based on observations of magnetic reconnection and merging in

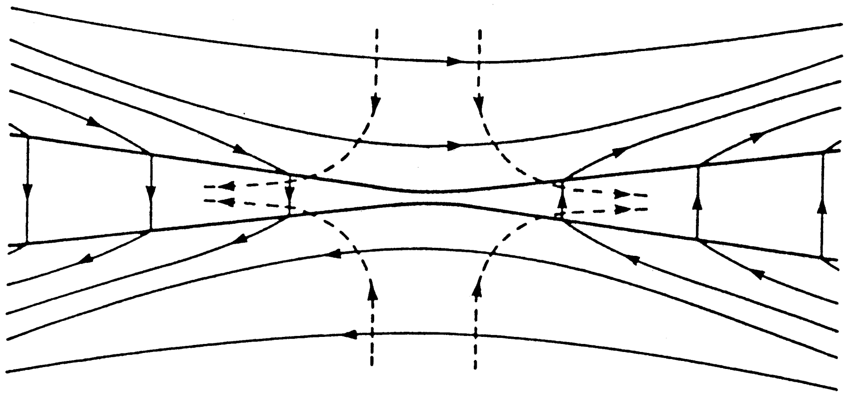

Fig. 2. Petschek's (1964) model for steady-state reconnection. The magnetic field is convected into a neutral line or plane where reconnection occurs heating the plasma and accelerating particles. The rate at which reconnection occurs is greatly enhanced over earlier models due to a much more effective removal of the heated plasma and reconnected field lines through slow mode shocks, depicted by the lines passing through the discontinuities in the directions of the magnetic field lines. The dashed lines indicate the direction of plasma flow. Other models for reconnection were reviewed by Vasyliunas (1975).

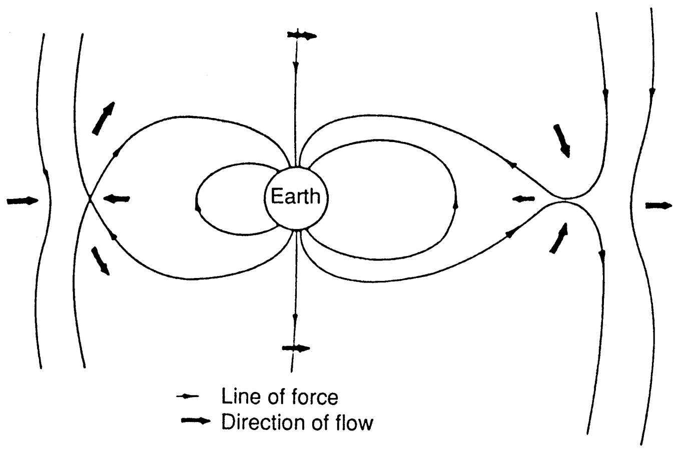

Fig. 3. Dungey's (1961) model for an open magnetosphere containing magnetic neutral points both ahead and behind the magnetosphere. Magnetic field lines joined to the earth reconnect with magnetic field lines in the solar wind at the neutral point on the solar side, magnetic flux is dragged from the solar to the antisolar side by the solar wind, reconnect at the neutral point on the antisolar side, and convect back to the solar side through the interior of the magnetosphere. Observational data suggest that reconnection on the solar side occurs rapidly, at a rate determined by the global requirements and not by the local plasma conditions: when the orientation of the magnetic field in the solar wind changes abruptly the reconnection rate adjusts immediately. A suggested implication for solar physics is that 'explosive' reconnection in solar flares is implausible, e.g., Akasofu (1984).

the terrestrial magnetosphere. As first pointed out by Dungey (1961) the magnetic field lines joined to the earth and those joined to the sun reconnect at two points on the solar and antisolar side of the earth, respectively, as illustrated in Fig. 3. It is thought that the rate of reconnection on the solar side is determined by the global structure of the magnetic field and not by the properties of the localised region where reconnection occurs, e.g., the reviews by Burch (1974) and Sonnerup (1979). These and related magnetospheric phenomena have been discussed by Akasofu (1977). Spicer (1982) and Akasofu (1984), amongst others, emphasised that reconnection is a driven process and should not be assumed as the driving mechanism for solar flares.

Syrovatskii (1966) proposed one specific model for a flare based on reconnection, cf. Anser (1973) however. Syrovatskii and Somov (1980), Syrovatskii (1981) reviewed such models, cf. also Somov and Titov (1985*a,b*) and Somov (1986*b*).

Acceleration of fast particles during reconnection is attributed to the electric field induced by the changing magnetic field. Particles in the current sheet separating the regions of oppositely directed magnetic field have orbits of a figure-of-eight character, e.g., Weiss and Wild (1964), Speiser (1965). Acceleration during reconnection was discussed, e.g., by Pikel'ner and Tsytovich (1967) and Bulanov and Sasarov (1976), cf. also Friedman and Hamberger (1968), Friedman (1969).

Magnetic annihilation models may be separated into two classes: reconnection models and tearing mode models. The theory of tearing modes was first developed by Furth, Killeen and Rosenbluth (1963), cf. also Pritchet, Lee and Drake (1980). The concept of tearing modes was incorporated into a flare model by Spicer (1977, 1981, 1982). In a model of magnetic energy dissipation based of tearing modes, the magnetic energy is released at a large number of localised sites associated with magnetic islands, as illustrated schematically in Fig. 4. This is to be compared with a reconnection model which has only one site (a current layer) where reconnection is occurring.

Acceleration by Parallel Electric Fields

The simplest conceivable mechanism for the acceleration of fast particles is by a parallel electric field $E_{\parallel}$. (A perpendicular electric field, except near a neutral plane in the magnetic field, causes an $\boldsymbol{E}\times\boldsymbol{B}$ drift for all particles and leads to no net acceleration.) There must be a large potential drop available to drive a solar flare, and a potential of order 10^{10} V was suggested, e.g., Swann (1933), Sweet (1958), Jacobsen and Carlqvist (1964), Alfvén and Carlqvist (1967), and Colgate (1978). A simple argument leading to this value is as follows. The power dissipated in a flare, which is up to 10^{22} W in a large flare, is electrodynamic and so one must be able to express it in terms of the rate work is done by a current against an electric field; power equals current times potential drop, and the available currents of up to 10^{12} A (e.g., Moreton and Severny 1968, Hagyard 1988) then imply a potential drop of order 10^{10} V. There are two arguments favouring a model based on a current of about 10^{12} A. One argument is that, to within a factor of order unity, vector magnetic field observations of active regions in which flares occur imply a

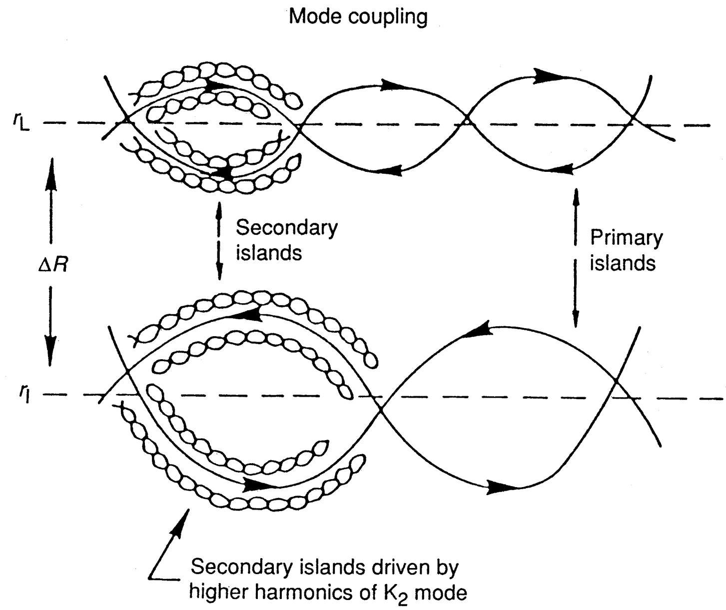

Fig. 4. The cartoon drawn by Spicer (1981) to illustrate how multiple regions of reconnection might result from coupling between tearing modes. A single reconnection region has a very small volume and this severely limits the rate at which magnetic field can be reconnected and plasma reprocessed. Many reconnection sites operating simultaneously are required for reconnection to occur fast enough to explain a flare.

maximum current of this order, and the regions of large current correlate with the positions of flare kernels, e.g., Lin and Gaizauskas (1987), Machado *et al.* (1988). A second argument is that a current of about 10^{12} A is the maximum that could flow into the corona because higher currents would self pinch. The self-pinching effect prevents a current flowing along a guiding magnetic field when the self-field due to the current is comparable with the guiding field. Consider an idealised model with a current confined to a cylinder. The field at a radius R due to the current I for $r < R$ is then $B = \mu_0 I/\pi R$. A typical strength for the guiding field in an active region of $B = 0{\cdot}15\,\mathrm{T}$, and a typical radius of the footpoints of flaring flux tubes might be $R = 3 \times 10^6$ m. In this case the self-field equals the guiding field for $I = 1 \times 10^{12}$ A. Larger currents can flow only either in stronger guiding fields ($B > 0{\cdot}15\,\mathrm{T}$) or in larger current channels ($R > 3 \times 10^6$ m). It follows that directed current flowing from one footpoint to the other (excluding the possibility of many oppositely directed currents through the corona) cannot exceed about 10^{12} A. It is assumed here that the current associated with a flare is such a directed current.

It follows that if the current is about 10^{12} A then the potential must be about 10^{10} V to account for the power in a flare. In the remainder of this review a potential of order of 10^{10} V is assumed, but two comments on this assumption are appropriate. The first comment is that there is no direct evidence for a potential of 10^{10} V; the arguments in favour of it are indirect. One might expect some evidence of acceleration producing particles with energy of order 10^{10} eV, but there is little such evidence. An indirect argument relates to the potential drop across a coronal mass ejectum (CME): if the velocity of the CME is attributed to an $\boldsymbol{E} \times \boldsymbol{B}$ drift then the implied E times the transverse dimension of the CME implies a potential of order 10^{10} V; the indirect argument is that such a potential must be available in the circuit in which the CME is modeled as a capacitor that takes off as it charges up. On the other hand, there is evidence that impulsive phase acceleration does favour electrons with energy of about 10^5 eV. The second comment relates to a possible alternative model in which the potential 10^5 V is taken as the starting point. Suppose one assumes that the available potential drop is 10^5 V and not 10^{10} V. Then to account for a power of 10^{22} W requires a current of 10^{17} A, which greatly exceeds the maximum possible direct current of about 10^{12} A that can flow through the corona. Hence a current of 10^{17} A would need to flow in at least 10^5 current channels, with half flowing up and half flowing down. Such a model has been suggested by Lin and Schwartz (1987). Although such a model is not discussed here, it is regarded as a possible alternative model that should be developed further. An obvious difficulty with such a model is how such a complicated current pattern could be set up.

One class of flare models is based on this potential collapsing to a localised double layer in the corona. Following Alfvén and Carlqvist (1967) this type of model is referred to by the name of *current interruption*, which is somewhat misleading (a preferable alternative is 'current disruption'). A better way of describing this type of model is one involving dissipation of a parallel current (e.g., Spicer 1982), which contrasts it with reconnection models that involve dissipation of a current flowing in a current sheet perpendicular to the mean field on either side of the sheet. The idea that double layers might be important in solar flares and other astrophysical contexts was discussed in reviews by, e.g., Alfvén (1977), Block (1978) and Raadu (1989). Double-layer like structures called electrostatic shocks are a common feature in the earth's auroral zones, where they are associated with fluxes of accelerated particles e.g., the review by Mozer *et al.* (1980). The acceleration of electrons by the parallel electric fields in the auroral zones of the earth is not adequately understood, although they are of considerable interest in connection with the generation of the auroral kilometric radiation (AKR). The most recent data from the *Viking* spacecraft suggest that the electrons that generate AKR are in a 'forbidden' region of phase space, which is accessible only to electrons experiencing a time-dependent parallel electric field (Louarn *et al.* 1990). The implications of this observation for analogous solar phenomena (microwave spike bursts) have yet to be considered.

Acceleration by a parallel electric field is usually treated in terms of a runaway process. Consider the equation of motion of an electron in the presence of an electric field and of a frictional force due to collisions with a collision frequency $\zeta_e(v) = \zeta_e(V_e/v)^3$, where V_e is the thermal speed of electrons:

$$\frac{d\boldsymbol{v}}{dt} = -\frac{e\boldsymbol{E}}{m} - \zeta_e \boldsymbol{v}\left(\frac{V_e}{v}\right)^3. \tag{2.13}$$

It follows that electrons with speed

$$\frac{v}{V_e} > \left(\frac{E_D}{E}\right)^{1/2}, \quad E_D = \frac{mV_e\zeta_e}{e}, \tag{2.14}$$

are freely accelerated. These are the *runaway* particles, and E_D is the *Dreicer* field (Dreicer 1959, 1960). Runaway acceleration was discussed in connection with solar flares by, e.g., Takakura (1971), Norman and Smith (1978), Kuijpers, van der Post and Slottje (1981), Tsuneta (1985) and Holman (1985*a,b*). In the presence of anomalous resistivity the collision-like process providing the frictional force in (2.13) is due to scattering by the waves that cause the anomalous resistivity, usually assumed to be ion sound waves (e.g., Takakura 1988) or lower hybrid waves.

There is a severe limitation on runaway acceleration that arises from the fact that the runaway electrons and ions constitute an electric current (Hoyng 1977*a,b*, Spicer 1983, Holman 1985*a*). In circuit language the current cannot change faster than on the inductive time scale which is the ratio of the inductance to the resistance. The inductive time scale depends on the details of the model and for a flaring region this time scale is at least as long as about the duration of the impulsive phase, e.g., Melrose and McClymont (1987). It follows that the number of particles (electrons or ions) that results from runaway acceleration cannot produce a current significantly different from that flowing before the onset of the flare. The current estimated from observation is of order 10^{12} A (e.g., Moreton and Severny 1968, Hagyard 1988), which corresponds to electrons flowing at the rate $10^{31}\ \mathrm{s}^{-1}$. Thus the runaway process alone cannot produce an electron flux of $\gtrsim 10^{35}\ \mathrm{s}^{-1}$ required to account for a solar flare.

3. Bulk Energisation of Electrons

Solar flares may be classified in several ways, e.g., Švestka (1976), who separated flares into 'low-temperature' and 'high-temperature' classes, Spicer and Brown (1980), who distinguished between $\boldsymbol{J}_\parallel$ and $\boldsymbol{J}_\perp$ driven energy release, and Bai and Sturrock (1989), who defined five classes of flares. Two of the classes defined by Bai and Sturrock (1989) are analogous to the two classes defined by Švestka (1976): these are their 'thermal hard X-ray flares' and 'nonthermal hard X-ray flares', and a third class is defined by an associated filament eruption. The other two classes are defined in terms of γ-ray and energetic particle signatures. It seems that in all solar flares a substantial fraction of the primary energy release is into electrons with energies 10 to 100 keV, which corresponds to a temperature of 10^8 to 10^9 K. For example, Duijveman, Hoyng and Machado (1982) estimated that > 20% of the dissipated power goes into such electrons. The mechanism involved in the so-called *bulk energisation* of the electrons is poorly understood. Some existing ideas were reviewed by Ramaty *et al.* (1980) and more recently by Benz (1987). In the discussion here emphasis is placed on an interpretation in terms of energy release in a weak double layer. However, some of the ideas are speculative.

Bulk Energisation as an Enhanced Form of Coronal Heating

The energy release mechanism in a solar flare is probably related to the coronal heating mechanism, e.g., Heyvaerts and Priest (1983), Ionson (1985*a,b*). A simple interpretation of the bulk energy release in a solar flare is that it is an extreme form of the heating of coronal flux tubes. Consider the heating of a coronal flux tube at various different rates. If the flux tube is heated slowly enough it reaches a steady state in which heating is balanced by conduction of energy back down to the chromosphere. The temperature of the flux tube is determined by the balance between the heating rate and the energy loss rate. In a given flux tube, increased heating should lead to an increase in the temperature. The energy spectrum of the heated electrons is maintained nearly Maxwellian provided that the collisional energy exchange time is shorter than other relevant times. Assuming that the heat goes primarily into electrons, it is redistributed to the ions through collisions on a time scale that is longer than the time for the electrons to thermalise by a factor ≈ 43 (e.g., Trubnikov 1965).

For sufficiently high temperatures this simple picture must change because conduction in the usual sense becomes ineffective when the collisional mean free path becomes comparable to the length of the flux tube. The energy transport must then be through either a collisionless conduction front (e.g., Brown, Melrose and Spicer 1979, Smith and Lilliquist 1979) or by freely propagating electrons (e.g., Emslie 1983). The column stopping distance for an electron with energy E is (e.g., de Jager *et al.* 1987)

$$D_e \approx \left(\frac{E}{mc^2}\right)^2 \frac{1}{2\pi r_0^2 \ln\Lambda}, \tag{3.1}$$

where $r_0 = e^2/4\pi\varepsilon_0 mc^2 \approx 2 \cdot 82 \times 10^{-15}\,\mathrm{m}$ is the classical radius of the electron and $\ln\Lambda \approx 20$ is the Coulomb logarithm. For a typical density $n = 10^{16}\,\mathrm{m}^{-3}$ the mean free path $\lambda_e = D_e/n$ is longer than the typical length $L = 10^7\,\mathrm{m}$ of a flux tube for $E/mc^2 \gtrsim 10^{-2}$ or $E \gtrsim 5\,\mathrm{keV}$.

A flare may be regarded as the result of rapid heating where the mean free path is greater than the length of the flux tube so that an enhanced form of energy transport away from the heated region is required.

There is evidence from both hard X-ray and microwave data that the energy release during the impulsive phase of a solar flare occurs in many localised regions on a time scale that can be $\lesssim 100\,\mathrm{ms}$. The location of the energy release is called a flare kernel; the cooling by conduction of impulsively heated kernels was discussed by Brown, Craig and Karpen (1978). The individual energy release episodes observed in X rays include elementary flare bursts (EFBs) of duration 5 to 15 s (de Jager and de Jonge 1978) and spikes of duration around 100 ms. This structure in both time and space of the energy release is interpreted in terms of energy release in individual flux tubes in a picture in which the magnetic structure consists of many intertwined flux tubes, e.g., Sturrock *et al.* (1984). Interpretation of the data on individual flare kernels leads to the conclusion that a range of parameters occurs in different flares. For example, Batchelor *et al.* (1985) found $n \approx 10^{15}\,\mathrm{m}^{-3}$, $B \approx 10^{-2}\,\mathrm{T}$, $L \approx 10^7\,\mathrm{m}$, and de Jager *et al.* (1987) found $n \approx 10^{17}\,\mathrm{m}^{-3}$, $B \approx 0{\cdot}14\,\mathrm{T}$, $L \approx 3{\cdot}5 \times 10^{5}\,\mathrm{m}$ and $T \approx 5 \times 10^{8}\,\mathrm{K} \approx 50\,\mathrm{keV}$.

In summary, the primary energy release in the impulsive phase of a solar flare appears to occur in many localised energy release events in which the electrons are heated up to a mean energy in the range 10 to 100 keV on a time scale that can be as short as 100 ms. This heating is too rapid to be balanced by collisional thermal conduction of the heat back to the photosphere.

The Need for Enhanced Resistivity

Most heating mechanisms for the solar corona and all specific processes suggested for bulk heating of solar flare plasma may be interpreted in terms of dissipation of electric currents. Thus, let us suppose that the local heating rate per unit volume is ηJ^2, where η is the electrical resistivity and $J = |\boldsymbol{J}|$ is the current density. It was argued by Melrose and McClymont (1987) that the classical value of the resistivity cannot account for energy release at the required rate in a flare and that it is essential that the resistivity be anomalous. Before summarising this argument it is relevant to note that not all bulk heating mechanisms can be described in terms of a resistivity. One exception is the mechanism for dissipation of Alfvén waves that involves phase mixing and so relies on the viscosity rather than the resistivity, e.g., Heyvaerts and Priest (1983). A more relevant exception applies when energy release occurs into mass motions rather than into random thermal motions, as implied by the concept of resistive dissipation. Nevertheless most energy release processes involving a current may be expressed in the form ηJ^2 with η identified as some equivalent resistivity. The following argument is that this equivalent resistivity must be enhanced above the Spitzer resistivity.

The local heating time τ_H may be defined as the ratio of ηJ^2 to the thermal energy density $n_e m_e V_e^2$ in the electrons:

$$\frac{1}{\tau_H} = \frac{\eta J^2}{n_e m_e V_e^2}. \tag{3.2}$$

The current density is related to the drift velocity v_D of the electrons relative to the ions by

$$J = -e n_e v_D. \tag{3.3}$$

The resistivity may be written in terms of an effective collision frequency ζ_{eff} between the electrons and the ions:

$$\eta = \frac{m_e \zeta_{\mathrm{eff}}}{n_e e^2} = \frac{\zeta_{\mathrm{eff}}}{\varepsilon_0 \omega_p^2}. \tag{3.4}$$

Then (3.2) implies

$$\frac{1}{\tau_H} = \zeta_{\mathrm{eff}} \frac{v_D^2}{V_e^2}. \tag{3.5}$$

It is now argued that if one supposes that the dissipation is due entirely to Coulomb interactions, then an inconsistency results. For Coulomb interactions, $\zeta_{\rm eff}$ is equal to the classical collision frequency ζ_0:

$$\zeta_0 = \frac{\omega_p \ln \Lambda}{4\pi n_e \lambda_{De}^3}, \tag{3.6}$$

where $\lambda_{De} = V_e/\omega_p$ is the Debye length. For the solar corona typical numerical values give $\zeta_0 \approx 3\times 10^{-4} n_e T_e^{-3/2}\,{\rm s}^{-1}$, and for $n_e = 10^{16}\,{\rm m}^{-3}$ and $T_e = 10^8\,{\rm K}$ this gives $\zeta_0 \approx 3\,{\rm s}^{-1}$. An important part of the present argument is that the drift velocity v_D cannot exceed the threshold speed for various drifting instabilities, because otherwise the resistivity would be anomalous, contrary to the present hypothesis that the resistivity is given by the collisional value. The appropriate threshold is thought to be of order the ion sound speed v_s, in which case the factor in (3.5) satisfies

$$\frac{v_D^2}{V_e^2} \lesssim \frac{v_s^2}{V_e^2} \approx \frac{m_e}{m_i}, \tag{3.7}$$

where m_i is the mass of the ions. Then (3.5) implies $\tau_H \gtrsim 500\,{\rm s}$. Now the heating time for the entire volume of the flare cannot be shorter than the local heating time, and hence classical collisional effects cannot account for the heating on time scales of the order of a second or less, as required to account for the time structures observed in flares, e.g., Holman (1985*b*).

The argument that an enhanced form of resistivity is needed is rather an argument that classical resistivity is inadequate. Although it appears that conventional forms of anomalous resistivity are adequate to account for the average dissipation in a flare, e.g., Melrose and McClymont (1987), it is doubtful that this also applies to the faster rate of dissipation in the fine structure in flares, specifically EFBs and spikes. An acceptable model may well require a faster rate of dissipation, perhaps through hyperresistivity (e.g., Strauss 1988) or through double layers, as discussed below.

Dissipation in Thin Current Channels

Let us now suppose that the dissipation is due to anomalous resistivity. (As indicated above, even anomalous resistivity may be inadequate to account for dissipation in EFBs and spikes, but at least for the average dissipation during the impulsive phase it seems adequate.) Then there are several constraints that have important implications. The following arguments are related to those given by Chiuderi (1983), Melrose and McClymont (1987) and Khan (1989). The arguments are presented only for ion sound turbulence; they apply for other forms of anomalous resistivity (e.g., Duijveman, Hoyng and Ionson 1981) but need modification in numerical details.

The assumption that the current is at the threshold for ion sound turbulence implies that the drift speed v_D in (3.3) is fixed at a value that differs from the ion sound speed $v_s \approx V_e/43$ by a factor of order unity, which is ignored here. Also the value of the resistivity (3.4) is determined to within a factor of order unity and corresponds to $\zeta_{\rm eff} \approx \omega_p v_s/V_e$. Thus for ion sound turbulence we take

$$J = n_e e v_s, \quad \zeta_{\rm eff} \approx \omega_p \frac{v_s}{V_e}. \tag{3.8}$$

The current density $J = n_e e v_s$ must be restricted to thin sheets or lines. The argument for this follows from the requirement that the magnetic field associated with the current not exceed the ambient field. Following Spicer and Brown (1980) let us discuss the cases of perpendicular and parallel currents separately. The case of a perpendicular current flowing uniformly in a sheet may be used to model a neutral sheet in the magnetic field. Let the sheet be of thickness $\ell_\perp$ and let B_0 and $-B_0$ be the value of the fields at the edges at $\ell_\perp/2$ on either side of the sheet. Then $\mathrm{curl}\,\boldsymbol{B} = \mu_0 \boldsymbol{J}$ implies $\ell_\perp = 2B_0/\mu_0 n_e e v_s$. The case of a parallel current is appropriate when considering, for example, dissipation due to double layers. For a parallel current (along the z axis) confined to a sheet of thickness $\ell_\perp$ (with normal along the x axis), the current can flow only if the self field $B_x = \pm\frac{1}{2}\mu_0 J \ell_\perp$ (at the edges of the sheet where it has its maximum value) is less than the confining field $B_z = B_0$. This implies $\ell_\perp < 2B_0/\mu_0 n_e e v_s$. For a current flowing uniformly in a cylinder of radius r, the current can flow only if the self field $B = \frac{1}{2}\mu_0 J r$ (at the edge of the cylinder) satisfies $r < 2B_0/\mu_0 n_e e v_s$. In all cases, the thickness $\ell_\perp$ of the region in which the current flows is restricted by

$$\ell_\perp \lesssim \frac{2B_0}{\mu_0 n_e e v_s} = 2\frac{v_A}{V_e}\frac{c}{\omega_p}. \tag{3.9}$$

The geometry of the channels in which the current flows is not known. Two opposite extreme assumptions are (a) the current flows in sheets of dimensions $L_\perp \ell_\perp$, where $L_\perp$ is the thickness of the flaring flux tube, and (b) the current flows in cylinders with radius $r = \ell_\perp$, where in both cases $\ell_\perp$ is determined by replacing the inequality by an equality in (3.9). Let there be N such current channels in which the dissipation is occurring at any given time. Then the filling factor f of the total volume $V = L_\parallel L_\perp^2$ that is filled by the current channels is given by

$$f = N\frac{\ell_\parallel}{L_\parallel}\left(\frac{\ell_\perp}{L_\perp}\right)^p, \tag{3.10}$$

with $p = 1$ for case (a) and $p = 2$ for case (b), and where $\ell_\parallel$ is the length along the the flux tube in which the energy release is occurring at any one time.

Consider a model for bulk energisation in which the energy released in a flare goes into heating all the electrons in the volume V. The average power released in the volume V is $P = f\eta J^2 V$. On dividing total energy released by P one obtains the characteristic time τ_F for the flare:

$$P = f\eta J^2 V = \frac{n_e \varepsilon_e V}{\tau_F}, \tag{3.11}$$

where $\varepsilon_e = m_e V_e^2$ is the mean energy of the heated electrons. On inserting the value (3.4) with (3.8) for the resistivity, one may solve (3.11) for f:

$$f = \frac{V_e}{v_s \omega_p \tau_F} \left(\frac{V_e}{v_s}\right)^2 . \tag{3.12}$$

With $\omega_p = 6 \times 10^9\,\mathrm{s}^{-1}$ and $\tau_F = 3 \times 10^2\,\mathrm{s}$, one estimates $f \approx 10^{-7}$. Thus *at any given time the energy release is occurring in a very small fraction of the volume of the flaring region.* The total number N of energy release regions is strongly dependent on their assumed geometry, and limits on it may be estimated from (3.10). On inserting numerical values in (3.9) one estimates that $\ell_\perp$ is of order of magnitude 1 m to 1 km, and assuming $L_\perp = 10^6\,\mathrm{m}$, one finds $\ell_\perp/L_\perp$ is of order 10^{-6}. Thus one estimates from (3.10) that N is of order $L_\parallel/10\ell_\parallel$ for $p = 1$ and of order $10^5 L_\parallel/\ell_\parallel$ for $p = 2$.

It might be remarked that the difficulties associated with including anomalous resistivity were pointed out by Chiuderi (1981), who suggested that the difficulty in generating and maintaining the conditions for anomalous resistivity favoured a model based on transient episodes of energy release. Chiuderi (1981) referred to 'numerous mini-reconnections, occurring randomly in space and time' and recalled that many reconnection sites are invoked in a model proposed by Levine (1974*a,b*). Chiuderi (1981) went on to say that 'a quantitative analysis of this mechanism is badly needed', but as yet no such analysis seems to have been presented. Similar ideas were advanced by Melrose and McClymont (1987) and Khan (1989), who envisaged dissipation of parallel current rather that the perpendicular currents involved in magnetic reconnection or tearing.

In summary, the energy release must involve an enhanced (over collisional) dissipation process and, because such dissipation is confined to localised regions, the energy release must occur in many localised regions. For all the electrons in the flaring region to be energised, they must all pass through one or more of these energy release regions during the flare. This requires that either

(i) there be a flow of plasma through each energy release region,
(ii) the energy release regions move through the plasma, or
(iii) the energy release regions form and break up many times and in many locations through the flaring region.

Further development of these ideas requires an explicit identification of the energy release mechanism and of the geometry of the energy release regions.

Reconnection Models

Reconnection is the favoured model for energy release in solar flares and is the favoured mechanism for bulk energisation of electrons. Most discussions of the energy release relate to coronal heating, on which there is an extensive literature, e.g., Levine (1974*a,b*), Van Hoven (1979), Mullan and Levine (1981), Duijveman, Hoyng and Ionson (1981), Galeev *et al.* (1981), Chiuderi (1981, 1983), Coppi (1983), Heyvaerts and Priest (1984), Steinolfson and Van Hoven (1984), Spicer, Mariska and Boris (1985), Tsuneta (1985), Martens, van den Oord and Hoyng (1985), Somov and Titov (1985*a,b*), Browning and Priest (1986), Chiueh and Zweibel (1987), Tajima *et al.* (1987), Hénoux and Somov (1987), Low and Wolson (1988), Mikic, Barnes and Schnack (1988) and the reviews by Kuperus, Ionson and Spicer (1981), Ionson (1985*a,b*) and Somov (1986*a*). For the purpose of discussion of reconnection models here, it is assumed that the energy release in a flare is an enhanced form of coronal heating. However, the details are ill defined, specifically concerning how the energy release occurs locally and how the global energy release is related to the local energy release. Some authors envisage a single large current sheet while others envisage several or a very large number of smaller current sheets. Some authors also argue in favour of runaway acceleration in an electric field E that is comparable to the Dreicer field E_D, cf. (2.14). Here it is argued that a single large current sheet is unacceptable, and that runaway acceleration is unacceptable as a bulk energisation process.

Dissipation in a single large current sheet seems implausible. One difficulty is that if such a large current sheet did form, then various instabilities should broaden it and induce turbulent motions in it (Uchida and Sakurai 1977, Chiueh and Zweibel 1987, Strauss 1988, Strauss and Otani 1988) so that effectively it would break into many smaller regions. A more serious difficulty concerns the flow of plasma through the sheet. An important constraint is that the mass inflow and outflow rates must balance. This requires

$$n_{\rm in}v_{\rm in}A_{\rm in} = n_{\rm out}v_{\rm out}A_{\rm out}, \tag{3.13}$$

where n, v and A denote, respectively, the number densities, flow speeds and areas over which the inflow and outflow occur. For simplicity let us ignore any compression of the plasma and set $n_{\rm in} = n_{\rm out}$. Then (3.13) implies $v_{\rm in}A_{\rm in} = v_{\rm out}A_{\rm out}$. The ratio $A_{\rm out}/A_{\rm in}$ is equal to $\ell_\perp/L_\perp$ or to $\ell_\perp/L_\parallel$, depending on whether the outflow is perpendicular or parallel, respectively, to the magnetic field outside the current sheet. The very small value of $\ell_\perp$ estimated above then implies that $v_{\rm in}/v_{\rm out} = \ell_\perp/L_\perp$ or $\ell_\perp/L_\parallel$ is very small. The outflow speed is restricted to about v_A, placing a severe constraint on $v_{\rm in}$ and hence on the rate at which plasma can be processed through a single large but thin current sheet. For example, to produce observed energetic electron fluxes between 10^{35} and $10^{36}\,{\rm s}^{-1}$ with V_e of order $0\cdot1c$ and $n_e = 10^{16}\,{\rm m}^{-3}$ over an area $\ell_\perp L_\parallel$ with $\ell_\perp \approx 1\,{\rm m}$ requires $L_\parallel \approx 10^{18}\,{\rm m}$, and $\ell_\perp \approx 1\,{\rm km}$ requires $L_\parallel \approx 10^{15}\,{\rm m}$, that is, a sheet of order between $0\cdot1$ and 100 light years long, which is absurd. It is obvious that either the current sheet is much thicker than the value $\ell_\perp$ estimated by (3.9), or that there are many thin sheets.

It is worth emphasising that the argument given above in estimating $\ell_\perp$ cannot be avoided in any simple way. A single current sheet with a current density at the threshold for the ion sound instability can be no more than 1 to 10^3 m thick in the corona. It is then not possible to have a large enough electron flux escape from such a current sheet to account for the electron fluxes observed in solar flares. For models that invoke a single current sheet to be acceptable, e.g., as envisaged by Chiueh and Zweibel (1987), they must be regarded as involving implicit assumptions on energy release in thin sheets embedded within the thick sheet.

It might be remarked that Loran and Brown (1985) used an argument related to the foregoing to estimate a minimum B in a reconnection site. They wrote the inflow speed v_M as ζv_A with ζ a parameter of order unity, and they wrote the outgoing electron flux as $\mathcal{F} \lesssim \zeta n_e V_e A$, with $\zeta \approx 0 \cdot 2$ from Duijveman, Hoyng and Machado (1982). This model was used in a slightly modified form by de Jager *et al.* (1987) in interpreting a specific flare. In estimating the area A of the flaring region Loran and Brown (1985) wrote, in the notation used here, $A = \lambda L_\perp^2$ with $\lambda \leq 1$. This type of model cannot be a single flux tube in the sense envisaged here; for example, $L_\perp = 10^6$ m implies that $\lambda = \ell_\perp / L_\perp$ is between 10^{-6} and 10^{-3}. The large current sheet envisaged in such a model must be composed of many thin sheets the sum of whose volumes is only a small fraction of the total volume of the thick sheet.

Now let us assume that the energy release occurs in a thin current sheet and consider limitations on the energy release process. In an acceptable flare model, the primary energy release mechanism for the electrons in a single current sheet cannot be runaway acceleration, as is sometimes assumed (e.g., Tsuneta 1985). The reason for this is given at the end of §2 above. Specifically, if the energy release were due predominantly to runaway acceleration, then the power released could be written as $\dot{N}$ times the mean energy per runaway electron, and the current would be $I = -e\dot{N}$; the former requires $\dot{N} \gtrsim 10^{36}\,\mathrm{s}^{-1}$ and the current limitation argument requires $\dot{N} \lesssim 10^{31}\,\mathrm{s}^{-1}$. In an acceptable model, on average (over the energy release site) the electrons must flow out of local acceleration regions in nearly equal numbers in opposite directions along the magnetic field.

An energy release mechanism that is consistent with the requirement on current limitation is one in which the energetic particles are produced in the reconnection region by the hot plasma being squirted out the separatrices at about v_A (e.g., Vasyliunas 1975). There is observational evidence that this occurs in the plasma sheet in the earth's magnetotail (Lin *et al.* 1977). However, as it is neutral plasma that is squirted out, and the kinetic energy of this motion is primarily in the ions, for this mechanism to produce the electron fluxes inferred in solar flares, the energy must be transferred from the ions to the electrons. This seems an implausible mechanism for the bulk energisation mechanism in solar flares.

In summary, the suggestion that the primary energy release in solar flares is substantially into bulk energisation of electrons due to magnetic reconnection encounters two major constraints. One constraint follows from the fact that the dissipation must be due to enhanced (above collisional) equivalent transport coefficients; the current sheet in which the dissipation occurs must then be very thin (1 to 10^3 m). A much thicker current sheet envisaged by many authors can be consistent with this requirement only if it composed of many thin sheets which fill only a small fraction of the total volume. The other constraint effectively eliminates runaway acceleration, at least in its simplest form, as the primary acceleration mechanism. In effect when these constraints are taken into account no existing model for bulk energisation due to magnetic reconnection is acceptable as the primary energy release mechanism in a flare.

Strong Double Layer Models

As already noted, in the terminology of Spicer and Brown (1980), reconnection models involve dissipation of a perpendicular current. An alternative consists of models that involve dissipation of a parallel current. Models based on dissipation of a parallel current are usually described in terms of electrostatic double layers or electrostatic shocks.

It was suggested by Jacobsen and Carlqvist (1964) that solar flares might be due to current interruption and this idea was further developed by Alfvén and Carlqvist (1967) into what is now called the current-interruption model of solar flares, cf. also Carlqvist (1969, 1972, 1979*a,b*), Hasan and ter Haar (1978), Joyce and Hubbard (1978), Goertz and Borovsky (1983), Borovsky (1983, 1988), Tapping (1987) and the literature cited in the reviews by Carlqvist (1972), Block (1978), Torvén (1979), Smith (1982*a*) and Raadu (1989). The current-interruption model as originally proposed involves a single *strong* double layer.

Double layers may be classified as *strong* or *weak*. A strong double layer involves a coherent structure with electrons and ions flowing into the double layer and being transmitted or reflected depending on their initial momenta. The potential drop $\phi_{\rm DL}$ across the double layer is such that the potential energy is much greater than the thermal energy of the particles, i.e., $\phi_{\rm DL} \gg m_e V_e^2/e$. In contrast, a weak double layer involves a potential drop that is usually assumed to be in the range $0 \cdot 1 m_e V_e^2/e \lesssim \phi_{\rm DL} \lesssim 10 m_e V_e^2/e$. There is a region of wave turbulence driven by a current instability, and the potential drop of the double layer is embedded in this region of wave turbulence.

Weak double layers may be classified in terms of the wave mode involved. The two wave modes of most interest are the ion sound mode and the ion cyclotron mode. Ion sound double layers were the first to be identified (Sato and Okuda 1980, 1981) and have been discussed by, e.g., Okuda and Ashour-Abdalla (1982), Smith (1982*a,b* 1985), Chanteur *et al.* (1983), Barnes, Hudson and Lotko (1985). The electrostatic shocks in the earth's auroral zones are identified as ion cylotron double layers, e.g., Kan (1975, 1982), Hudson, Lysak and Mozer (1978), Goertz (1979), Temerin *et al.* (1982), Temerin and Mozer (1984), Smith (1986*a,b*) and Boström *et al.* (1988). From the point of view of acceleration of particles, a model involving a single strong double layer encounters difficulties that are similar to but more severe than encountered with a single current sheet. The difficulties include the following:–

(i) If the potential difference of $\phi_0 \approx 10^{10}\,$V required to account for the power in a large flare were to cross a single strong double layer, then one would expect particles with energy of order $10^{10}\,$eV to be accelerated directly. There is no evidence that this occurs.

(ii) The restriction on the number of particles accelerated to $\lesssim 10^{31}\,{\rm s}^{-1}$ is inconsistent with the fluxes of electrons inferred from hard X-ray bursts and observed in type III events (e.g., Lin 1985).

(iii) There is a conceptual problem as to how a single strong double layer might form across a flux tube with the dimensions of the energy release region in a flare; there is no model for how information might be communicated across the field lines to allow such a double layer to form.

(iv) The arguments given above concerning dissipation favour dissipation in many small localised regions.

In view of these difficulties a model based on a single strong double layer is not considered further here.

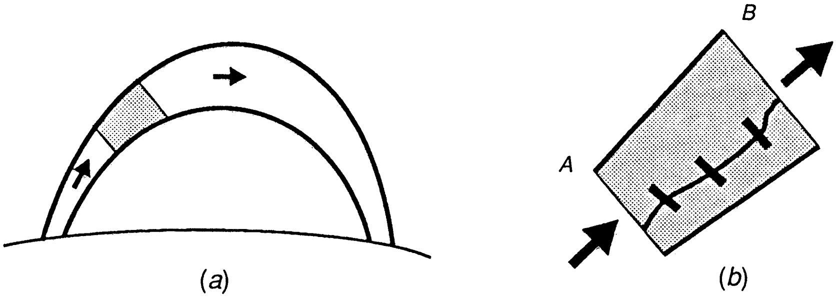

Fig. 5. A model for a solar flare based on multiple weak double layers is illustrated. (*a*) The electric current, denoted by the solid arrows, flows from one footpoint of a solar flux tube to the other through the corona; the energy release occurs in a flare kernel, indicated by the shaded region. (*b*) The available potential, assumed to be approximately 10^{10} V, develops across the flare kernel region, between the points denoted *A* and *B*, due to the shaded region becoming highly resistive. The current breaks up into very many filaments, only one of which is illustrated schematically. Very many multiple weak double layers form along each current filament, and three such double layers are indicated schematically by the dark boxes.

A Weak Double Layer Model

A more favourable model involving dissipation of a parallel current is one that involves weak multiple double layers. The idea is that the potential drop along one field line forms across several weak double layers, e.g., Hershkowitz (1981, 1985), Smith (1982*a*, 1985), Chan and Hershkowitz (1982), Bailey and Hershkowitz (1988). A specific model for solar flares based on multiple double layers was presented by Khan (1989). This model is illustrated in Fig. 5.

Khan (1989) assumed that each double layer has dimensions, cf. Barnes, Hudson and Lotko (1985),

$$\ell_{\parallel} \approx 10\lambda_D, \quad \ell_{\perp} \approx \left[\lambda_D^2 + \rho_s^2\right]^{1/2}, \quad \phi_{DL} = \xi m_e V_e^2/e, \tag{3.14}$$

where $\lambda_D = V_e/\omega_p$ is the Debye length, $\rho_s = v_s/\Omega_i$ is the ion-sound gyroradius, with $\Omega_i = q_i B/m_i$ the ion gyrofrequency, where ϕ_{DL} is the potential drop across the double layer and ξ is a parameter of order unity. The current is assumed to be confined to current channels with the current density in each channel at the threshold value

$$J_{DL} = n_e e v_s \tag{3.15}$$

for the generation of ion sound turbulence. Assuming a circular cross section for each channel, the current flowing in each channel is

$$I_{DL} = n_e e v_s A_{\perp}, \quad A_{\perp} = \pi \ell_{\perp}^2/4. \tag{3.16}$$

The power dissipated in each double layer is

$$P_{DL} = I_{DL}\phi_{DL} = \frac{\pi\xi}{4}(m_e n_e V_e^2)\ell_{\perp}^2 v_s. \tag{3.17}$$

The heating rate is proportional to the thermal energy density, and if a heating time τ_H is defined for each double layer, then one has

$$P_{DL} = \frac{m_e n_e V_e^2 \ell_\| A_\perp}{\tau_H}, \quad \tau_H = \frac{4}{\pi\xi}\frac{10(m_i/m_e)^{1/2}}{\omega_p}. \tag{3.18}$$

There are three constraints that need to be satisfied for such a model to be viable:–

(1) The number of double layers $N_\|$ along a given current channel must be such that $N_\|\phi_{DL}$ is equal to the total potential drop of order $\phi_0 \approx 10^{10}$ V.
(2) The number of current channels $N_\perp$ must be such that $N_\perp I_{DL}$ is equal to the total current, which is assumed to be of order $I_0 \approx 10^{12}$ A flowing through the flaring region.
(3) The total power dissipated $N_\| N_\perp P_{DL}$ must be equal to the total power of order up to 10^{22} W released in a flare.

Khan (1989) applied this model to the parameters for the flare discussed by de Jager *et al.* (1987) and found agreement for $N_\| \approx 1 \cdot 6 \times 10^5/\xi$ and $N_\perp \approx 5 \times 10^5$.

Speculations on the Energetics of Weak Multiple Double Layers

In attempting to formulate a model for the energy dissipation based on many short-lived double layers, a major difficulty is the absence of detailed discussion in the literature of the energetics of weak double layers. In order to proceed it is necessary to have some theory for the energetics. While it is inappropriate to attempt to formulate a detailed theory here, the following remarks outline some speculative ideas on how such a theory might be formulated.

The basic idea adopted here is that a discussion of the energetics of weak double layers could be based on a circuit approach. If one could identify the circuit parameters for a double layer then one could discuss its energetics. However, the circuit parameters have not been identified satisfactorily, cf. Sato and Okuda (1980), Smith (1986*a*). Nevertheless some progress can be made simply by examining plausible forms of the circuit parameters.

In the speculative discussion here it is assumed that the life of a double layer can be separated into three phases: a development phase, a main phase and a break-up phase.

Consider a localised region where a double layer forms and then breaks up. This region is part of a large circuit which has a total potential $\approx 10^{10}$ V greatly in excess of the maximum potential ϕ_{DL} that ever develops across the localised region. Before the current instability develops and generates ion sound turbulence, the potential drop ϕ across the localised region is negligible. The resistance R and the inverse capacitance $1/C$ are also negligible. Once the instability starts to develop, all these parameters become functions of time. The circuit equation may be written in the form

$$\phi(t) = R(t)I(t) + \frac{Q(t)}{C(t)}. \tag{3.19}$$

The length $\ell_{\parallel}(t)$ is that of the region in which the ion sound turbulence is present. This region is expected to expand from an initial small value to a maximum value; thus one expects $R(t)$ and $1/C(t)$ to change in time proportional to $\ell_{\parallel}(t)$ in accord with the familiar formulas

$$R(t) = \frac{\eta \ell_{\parallel}(t)}{A_{\perp}}, \quad \frac{1}{C(t)} = \frac{\ell_{\parallel}(t)}{\varepsilon_0 A_{\perp}} \tag{3.20}$$

for the circuit parameters, where for simplicity η and $A_{\perp}$ are assumed constant. As the double layer forms the current flowing through the region is partly diverted into building up the localised charge separation across the turbulent region leading to an increase in $Q(t)$. Energy must flow into the double layer in the form of a Poynting flux during this development phase but the details of how this occurs are uncertain.

There is one model for the energy flow into a double layer which is assumed to be created instantaneously (Carlqvist 1979, Raadu 1989). An important implication of the Carlqvist-Raadu model is that the appearance of a double layer in a circuit causes the current to be partially deflected, presumably around the double layer. This requires a time dependence of $I(t)$ in the localised region to describe the reduction in the current through the developing double layer as the current is deflected around the developing region.

It is impractical to attempt to discuss the energetics of this developing phase in a formal way for two reasons: first, the circuit equation has no energy integral when the capacitance depends on time and, second, the details of the energy flow into the region are not adequately understood.

In the main phase, once the double layer has formed it may be described by the five parameters ϕ_{DL}, R_{DL}, C_{DL}, I_{DL} and Q_{DL}, all of which may be considered constants in the main phase. The energy stored in the double layer is

$$E_{DL} = \frac{Q_{DL}^2}{2C_{DL}}. \tag{3.21}$$

An important point is that the power dissipated $P_{DL} = I_{DL}\phi_{DL}$, as given by (3.17) and (3.18), is much greater than the resistive loss $R_{DL}I_{DL}^2$ due to dissipation of the current in the resistance R_{DL} of the double layer. To see this, note that the ratio of the two powers may be evaluted using (3.20) with the resistivity given by (3.4) with (3.8), and with $I_{DL} = n_e e v_s A_{\perp}$ and $\phi_{DL} = \xi m_e V_e^2/e$:

$$\frac{I_{DL}\phi_{DL}}{R_{DL}I_{DL}^2} = \frac{\xi V_e^2}{v_s \ell_{\parallel} \zeta_{\text{eff}}} = \xi \frac{V_e^2}{v_s^2} \frac{\lambda_D}{\ell_{\parallel}}. \tag{3.22}$$

For plausible parameters this ratio is much greater than unity. It follows that most of the potential is across the capacitance, with

$$\phi_{DL} \approx Q_{DL}/C_{DL} \gg R_{DL}I_{DL}. \tag{3.23}$$

This suggests that the power going into the double layer should appear primarily as kinetic energy rather than as heat inside the double layer.

This raises the important question as to how directed acceleration in a double layer can lead to nearly equal fluxes of electrons (and ions, which we ignore for simplicity) escaping in opposite directions. The power dissipated in the double layer appears primarily as kinetic energy in electrons, and the question is how this results in electrons escaping from either end of the double layer region. One idea is to attribute this escape to diffusion of electrons through the double layer due to scattering by the ion sound turbulence. The idea is that the dissipation in a weak double layer is more analogous to ohmic dissipation in a metal than to dissipation in a strong double layer in that, as a result of the scattering by ion sound waves, the energy may be regarded as going primarily into 'heat' rather than into directed motion. The energy escapes from the double layer by the 'heated' electrons diffusing out both sides, rather than by directed flows of particles of opposite charges out opposite sides of the double layer.

The spatial diffusion coefficient is

$$D = \frac{V_e^2}{\zeta_{\rm eff}}. \tag{3.23}$$

Electrons diffuse a distance $\delta\ell$ in a time δt satisfying $D = (\delta\ell)^2/\delta t$. The escape time for a typical electron in the double layer is therefore

$$t_{\rm esc} \approx \frac{\ell_\parallel^2}{D}. \tag{3.24}$$

The power in the escaping electrons is of order

$$P_{\rm esc} = \frac{n_e m_e V_e^2 A_\perp \ell_\parallel}{t_{\rm esc}} = \frac{n_e m_e V_e^4 A_\perp}{\ell_\parallel \zeta_{\rm eff}}. \tag{3.25}$$

On comparing this to $I_{DL}\phi_{DL}$, one finds

$$\frac{P_{\rm esc}}{I_{DL}\phi_{DL}} = \frac{1}{\xi}\frac{\lambda_D}{\ell_\parallel}\left(\frac{V_e}{v_s}\right)^2. \tag{3.26}$$

If the power dissipated is to be identified with the loss of energetic electrons, then this ratio needs to be of order unity. With the estimates made in the literature on weak double layers cited above, this is possible only if $\ell_\parallel$ is considerably larger than that given by (3.14). This may be reasonable if the region of wave turbulence is much larger than the double layer itself.

Note that the electrons diffuse out both sides of the double layer region and do not flow in one side and out the other. Then the rate per unit time at which energetic electrons escape is not restricted by the current limitation presented above.

In the final phase envisaged here the double layer breaks up. The lifetime of the double layer may be estimated by regarding the double layer as a leaky capacitance. Its lifetime is then the capacitive timescale

$$t_{\rm life} = RC = \varepsilon_0 \eta = \frac{1}{\omega_p} \frac{v_s}{V_e}, \tag{3.27}$$

where (3.20) is used and then (3.4) and (3.8) are used. This estimate is somewhat longer than the lifetime indicated by simulation experiments, e.g., Sato and Okuda (1980). Perhaps the region in which the wave turbulence is excited is larger than the double layer itself, and $\ell_\parallel$ in (3.27) should be replaced by the longer length characteristic of this region. A constraint on a diffusion model is that the diffusive speed, which is the flow speed $v_{\rm esc}$ of the escaping electrons, must not exceed the thermal speed V_e of the electrons. One has

$$v_{\rm esc} = \frac{V_e^2}{\ell_\parallel \zeta_{\rm eff}} = V_e \frac{\lambda_D}{\ell_\parallel} \frac{V_e}{v_s}, \tag{3.28}$$

where (3.8) is used. With the parameters in (3.14) this escape speed is too large. This problem is also alleviated by assuming that the region in which the wave turbulence is excited is $\gg \ell_\parallel$. It may be concluded that a model in which the electrons diffuse out of the double layer due to scattering by the wave turbulence can account for the escape of the electrons in both directions from the double layer, but only if the region of wave turbulence is considerably larger than the double layer itself.

These somewhat speculative arguments may be summarised as follows:–

(1) The power input into a weak double layer is much greater than the power dissipated due to anomalous resistivity; from a circuit viewpoint, the power goes primarily into stored capacitive energy; from the viewpoint of dissipation, a weak double layer should be interpreted in terms of a hyperresistivity (e.g., Strauss 1988) rather than anomalous resistivity.

(2) The energy probably escapes from the double layer in the form of energetic particles diffusing out either end of the double layer through the effect of scattering by wave turbulence, but this model is acceptable only if the diffusive region is considerably larger than the double layer itself.

(3) The lifetime of the double layer may be interpreted in terms of a leaky capacitor model, subject to the same proviso as in (2).

(4) Following Carlqvist (1979) and Raadu (1989) one expects the development of a double layer to cause a deflection of current to paths outside the double layer; the implications of this have yet to be included in any model.

A Model for Bulk Energisation involving Many Short-lived Dissipation Regions

The foregoing discussion suggests that bulk energisation can occur due to electrons diffusing out of many weak double layers. Each double layer has a relatively short lifetime, but the average number and distribution of double layers remains constant. The power released per double layer may be written

$$I_{DL}\phi_{DL} = \omega_p\,(\lambda_D \ell_\parallel)^2\, n_e m_e V_e^2 A_\perp \ell_\parallel. \tag{3.29}$$

The fraction f of the total flaring volume filled by the double layers is given by

$$f = N_\perp N_\parallel \frac{A_\perp \ell_\parallel}{L_\perp^2 L_\parallel}, \tag{3.30}$$

where $N_\parallel$ is the number of double layers along a particular current channel and where $N_\perp$ is the number of current channels. The mean heating rate in the volume $L_\perp^2 L_\parallel$ may be identified from (3.29) by noting that $n_e m_e V_e^2 A_\perp \ell_\parallel$ is the thermal energy within each double layer, and hence the mean heating rate is f times $I_{DL}\phi_{DL}$ divided by this thermal energy. This give a heating time (the inverse of the heating rate)

$$\tau_H = \left[f\omega_p \left(\frac{\lambda_D}{\ell_\parallel}\right)^2 \right]^{-1}. \tag{3.31}$$

Now in (3.30) we assume $N_\perp = I_0/n_e e v_s A_\perp$ and $N_\parallel = e\phi_0/\xi m_e V_e^2$, where $I_0 \approx 10^{12}$ A is the total current and $\phi_0 \approx 10^{10}$ V is the total potential drop across the flaring region. As required (3.31) corresponds to

$$I_{DL}\phi_{DL} = \frac{n_e m_e V_e^2 L_\perp^2 L_\parallel}{\tau_H}. \tag{3.32}$$

The total number of double layers estimated by Khan (1989) is $N_\perp N_\parallel \approx 8\times 10^{10}/\xi$, implying $f = \xi/8\times 10^{10}$. One then finds that τ_H as given by (3.31) is of the appropriate order to apply to the time structures in the impulsive energy release in flares.

Acceleration of Runaway Electrons by Weak Double Layers

The foregoing outline of a model for bulk heating due to multiple weak double layers provides a possible alternative to reconnection as the basic energy release mechanism in flares. Besides further refinement of the model, there are two other features that need to be explored. These are runaway acceleration and the statistical acceleration of fast particles.

The acceleration of a nonthermal component of electrons should occur, in any model involving localised regions with anomalous resistivity, due to the runaway effect, e.g., as discussed by Norman and Smith (1978) and Holman (1985*a*). Briefly, following Holman (1985*a*) with the Coulomb collision frequency replaced by the effective collision frequency, (i) the critical speed for runaway in the presence of an electric field $E = m_e \zeta_{\rm eff} v_D/e$ is

$$v_c = \frac{m_e V_e^3 \zeta_{\rm eff}}{eE} = \left(\frac{V_e}{v_D}\right)^{1/2} V_e, \tag{3.34}$$

(ii) the fraction n_{er}/n_e of the electrons in the runaway regime is (Kaplan and

Tsytovich 1973, p. 134; Papadopoulos 1977, 1985)

$$\frac{n_{er}}{n_e} \approx \exp[-v_c^2/2V_e^2], \tag{3.35}$$

(iii) the time for formation of the runaway tail is

$$\tau_r = \left[\zeta_{\text{eff}}(V_e/v_c)^3\right]^{-1}, \tag{3.36}$$

(iv) and the rate per unit time that runaway electrons are produced is (Kruskal and Bernstein 1964, Knoepfel and Spong 1979)

$$\gamma \approx 0\cdot 35\,\zeta_{\text{eff}}\left(\frac{v_c}{V_e}\right)^{3/4} \exp\left[-2^{1/2}\left(\frac{v_c}{V_e}\right) - \frac{1}{4}\left(\frac{v_c}{V_e}\right)^2\right]. \tag{3.37}$$

On incorporating these ideas, one might hope to develop a theory for both bulk energisation and for the production of a nonthermal component on the basis of Khan's (1989) model of continually breaking up and reforming weak double layers.

The third feature that should be explored is the statistical acceleration of high energy particles as they randomly encounter the electric fields in the double layers. The statistical theory of acceleration due to such randomly distributed electric fields does not seem to have been developed. However, it was shown by Melrose and Cramer (1989) that the distribution function of fast particles encountering randomly distributed localised regions of Langmuir turbulence evolves in the same way as if the Langmuir waves were distributed uniformly with the same mean energy density, and similar methods may be applied to treat the case of randomly distributed weak double layers. Thus one expects the acceleration to be similar in form to that by Langmuir turbulence. The possible importance of such acceleration needs to be explored.

4. Prompt Acceleration of Ions

In the original separation of acceleration of solar flare particles into first and second phases (Wild, Smerd and Weiss 1963, de Jager 1969), it was envisaged that only electrons are accelerated in the first phase and that ions and relativistic electrons are accelerated in the second phase. However, this simple separation has been undermined by subsequent observations. Smerd (1975) pointed out that there is evidence for first-phase type acceleration of 10 to 100 MeV prompt protons in 'micro' proton events in the interplanetary medium, e.g., McDonald and Van Hollebeke (1985). The discovery of γ-ray line emission associated with flares (Chupp *et al.* 1973) implied the acceleration of ions with energies in excess of several tens of MeV accelerated in flares, and observations of these γ rays with a sufficiently short time scale showed (Chupp *et al.* 1981, Chupp 1983, Kane *et al.* 1986) that these ions are accelerated at the same time as impulsive electrons to within about a second, e.g., the review by Hudson (1985). This implies that what would earlier have been called a second phase acceleration mechanism operates on the timescale of about a second.

Several specific mechanisms for the prompt acceleration of such ions were discussed by, e.g., Melrose (1983), Bai and Dennis (1985), Ellison and Ramaty (1985), Decker and Vlahos (1986), Dröge and Schlickeiser (1986), Ohsawa and Sakai (1988), Smith and Brecht (1988, 1989*a,b*). The most notable result of these various investigations is that a variety of different acceleration mechanisms can be effective on the observed short time scale.

In summary, there has been a general swing of opinion away from the view that prompt acceleration of ions presents a major theoretical problem to the view that there are several possible ways of accounting for the data on prompt acceleration of ions. However, what the specific acceleration mechanism is and how it relates to bulk acceleration of electrons remains unclear.

5. Second Phase Acceleration

The original 'second phase acceleration' envisaged by Wild, Smerd and Weiss (1963) involves the acceleration of energetic ions and relativistic electrons starting at least several minutes after a flare. The second phase particles were assumed to be those that produce certain radio phenomena, notably type IV emission, and those that ultimately escape to be observed as solar cosmic rays. Acceleration by shock waves was the most favoured mechanism. More recent ideas have undermined the description of acceleration as occurring in two phases, and the concept of a second phase of acceleration is no longer useful, at least without considerable qualification. The more recent developments include the evidence that there is no clear separation between the time scales for prompt acceleration of electrons and of ions, and a re-examination of radio data that suggests more than two phases of acceleration.

These changes in emphasis have led to semantic difficulties with the use of 'second phase acceleration'. There are two different possible meanings. One relates to two (or more) separate phases of acceleration that involve separate populations of particles. The other relates to two stages in the acceleration of a single population of particles. To avoid this difficulty it is helpful to follow Bai *et al.* (1983) and to refer to the second stage in the latter case as 'second step' acceleration. Most acceleration mechanisms that are considered favourable for second phase acceleration, as originally envisaged, require a seed population of energetic particles which are then further accelerated. The two steps are the production of this seed population and the second step of their further acceleration.

Radio Data on Phases of Electron Acceleration

The radio evidence on acceleration was summarised by Smerd (1975) and Melrose and Dulk (1987). Smerd identified six phases implied by the radio data:–

(1) The classic first phase mechanism accelerates the electrons that produce type III bursts. The same population of electrons is believed to produce hard X-ray bursts, and a variety of other secondary observational signatures associated with their precipitation into the chromosphere. Smerd (1975) pointed out that if ions are accelerated to the same Lorentz factor they have energies in the range 10 to 100 MeV per nucleon, and it is in this context that he cited data on 'micro' proton events, as mentioned above.

(2) Electrons accelerated at shock fronts produce type II bursts. It is now thought that similar acceleration occurs at shock fronts in the interplanetary medium and leads to so-called SA events (Cane *et al.* 1981).

(3) Flare continuum radiation (e.g., Robinson 1985) appears following the impulsive phase. There are two distinguishable types of flare continuum, one (FCE) limited to the impulsive phase and the other (FCII) that extends longer and is associated with a type II burst. Both of these are distinct from type V bursts, which also involve a flare-related continuum and may be regarded as an extreme variant of type III emission (e.g., Suzuki and Dulk 1985).

(4) Moving type IV bursts (e.g., Stewart 1985) are attributed to plasmoids (plasma, magnetic field and energetic particles in a more or less self-contained configuration) ejected from the sun at $\approx 100\,\mathrm{km\,s^{-1}}$. The acceleration mechanism is not known.

(5) There is a prolonged phase of radio continuum ('late type IV' continuum) that lasts for several hours. This continuum is distinct from the flare continuum in its duration, its relation in time to the flare, and in the frequency range involved. However, it may be the same phenomenon as the type I continuum.

(6) Storms are not related directly to flares; they occur in active regions and can persist from hours to days. Storms can involve type I bursts, type I continuum, storm type III bursts and a variety of lesser known radio phenomena (e.g., Kai, Melrose and Suzuki 1985, Suzuki and Dulk 1985). As argued by Smerd (1975), the energetic electrons involved in storm emission cannot be attributed to trapped particles accelerated during the flare.

These various phases, and further radio phenomena discussed by Melrose and Dulk (1987), imply that acceleration of electrons occurs in a variety of contexts, and is not restricted to flares and to shock fronts.

The implications of radio data on acceleration are summarised in Table 1. The second column specifies the emission mechanism as plasma emission, gyrosynchrotron emission or electron cyclotron emission. The third column specifies the number of particles involved. For gyrosynchrotron emission this may be calculated from the radio data. For plasma emission and electron cyclotron maser emission one can have little confidence in any such calculation; the estimates are based on inferences from the number of particles observed in type III streams in the interplanetary medium. Further details on the electron distribution function are summarised in the remaining columns.

Acceleration Mechanisms for Radio Emitting Electrons

A variety of acceleration mechanisms is required to account for these radio data. Flare associated type III bursts are attributed to a small fraction of the electrons produced in the bulk energisation process escaping along open magnetic field lines. How the electrons find their way onto open field lines is not well understood. A possible explanation is simply that the magnetic configuration is partly stochastic in the sense that a fraction of field lines is open. Stochasticity (a 'random walk') of the field lines results from the presence of a spectrum of MHD turbulence and may be interpretated as causing

Table 1. Particle distribution functions for various types of solar radio bursts (Melrose and Dulk 1987)

Burst	Radiation mechanism	Total number of particles	Energy range, energy distribution	Pitch angle distribution	Remarks
1. First (flash or impulsive) phase					
Type III at $R \gtrsim 1\cdot 1 R_\odot$	Plasma	≈ 10 bunches $10^{31}-10^{33}$ per bunch	< 5–25 keV, bump-in-tail	Forward cone of angle $\approx 20°$	Faster e^- outpace slower e^-
Microwave impulsive	Gyrosynchrotron	$10^{14}-10^{17} > 10$ keV	10 keV–1 MeV, power law with $\gamma = 3$ to 5, or	Nearly isotropic	Most emission from e^- of $E \gtrsim 100$ keV
	Gyrosynchrotron	$\approx 10^{33}-10^{37}$	10^8-10^9 K Maxwellian	Nearly isotropic	Most emission from e^- of $E \gtrsim 100$ keV
Microwave spike					
Spike bursts	Cyclotron maser	?	10–100 keV	Loss cone	e^- in coronal loops
Type III-like	Plasma	?	10–100 keV	Forward cone	precursor to type III?
Flare continuum (FCE)	Plasma	$\gtrsim 10^{33} \gtrsim 10$ keV	Plateau or gap	Loss cone	e^- trapped in high coronal loops
2. Second phase					
Type II	Plasma	?	Plateau or bump-in-tail	Anisotropic	Shock associated
Microwave IV	Gyrosynchrotron	$10^{31}-10^{33}$	$0\cdot 5$–5 MeV, power law with $\gamma = 2$ to 3	Nearly isotropic	e^- trapped in low coronal loops
Moving type IV					
(a) Advancing front	Gyrosynchrotron	$\approx 10^{33} > 1$ MeV	1–3 MeV, power law with $\gamma \approx 5$ to 10	Nearly isotropic	Rare. Moves outward with shock
(b) Isolated source	Gyrosynchrotron?	$\approx 10^{33} > 0\cdot 1$ MeV	$0\cdot 1$–1 MeV, power law with $\gamma \gtrsim 4$	Nearly isotropic?	Plasmoids with $B \approx 0\cdot 3$–1×10^{-3} T
	Plasma?	$\lesssim 10^{34}$	≈ 10 keV, plateau or gap	Forward/loss cone	Early stages
Flare continuum (FCII)	Plasma	$\gtrsim 10^{34}$	$\gtrsim 10$ keV, plateau or gap	Loss cone	May be a continuation of FCE
3. Non-flare-associated					
Type I bursts	Plasma	?	$\gtrsim 3$–10 keV, gap?	Loss cone?	e^- injected into closed loops
Type I continuum	Plasma	?	$\gtrsim 3$–10 keV, gap?	Loss cone?	e^- trapped in high loops
Storm type III	Plasma	$\approx 10^4$ bunches per day $10^{31}-10^{33}$ per bunch	$\gtrsim 1$–10 keV, bump-in-tail	Forward cone	Similar to flare type III
Drift pairs	Plasma	?	?	Forward cone	type III-like?
S bursts	Plasma	?	?	Forward cone	type III-like?

a form of spatial diffusion across the mean field lines, e.g., Jokipii (1966). It is possible that the type III electrons are accelerated in a secondary process involving the absorption of electron cyclotron maser emission (Sprangle and Vlahos 1983). The impulsive microwave bursts are also attributed to the same initial population of bulk energised electrons, with the higher energy electrons dominating in this emission process.

Radio spike bursts, e.g., the review by Benz (1986), are regarded as radio signatures of the precipitating electrons. These may be separated into two classes, primarily on whether they are drifting rapidly in frequency or not. The favoured emission mechanism for the non-rapidly-drifting spike bursts is electron cyclotron maser emission (Holman, Eichler and Kundu 1980, Melrose and Dulk 1982), which is a topic of active current interest. Even for these bursts, plasma emission mechanisms (e.g., Vlahos, Sharma and Papadopoulos 1983) cannot be ruled out. An important detail in the electron cyclotron maser theory is the form of the free energy to drive the maser. Ideas on this free energy have been influenced by data on the precipitating electrons in the earth's auroral zone that produce the auroral kilometric radiation through electron cyclotron maser emission. The details of the generation and propagation of the electron beam are crucial to the detailed understanding of the way the maser is driven. However, the theory has other uncertainties when applied to the interpretation of radio spikes. An ongoing problem concerns the difficulty in accounting for the escape of radiation generated at the fundamental cyclotron frequency through the second harmonic absorption layer (Melrose and Dulk 1982). The interpretation is sufficiently uncertain that data on these spike bursts provide little useful information on the acceleration process. The other class of microwave spike bursts is type III-like in character (Elgarøy 1977). It is thought that these bursts, which often drift from low to high frequency, are signatures of the electrons escaping from the acceleration region. However, this idea has not been explored in sufficient detail to provide further insight into the acceleration process.

Type II bursts sometimes show a *herringbone structure* that is attributed to electrons streams propagating away from the shock front both ahead and behind the shock. There is direct evidence for such acceleration from *in situ* observations in the interplanetary medium, e.g., Potter (1981). A plausible acceleration mechanism is shock drift acceleration, e.g., Holman and Pesses (1983). However, it is not clear how this mechanism can lead to electron streams propagating away from the shock both ahead and behind the shock, as the radio observations on herringbone structure suggests, although the mechanism discussed by Cairns (1987) in a related context offers a basis for a possible explanation.

Type II bursts often involve a backbone that does not appear to be associated with electron streams. There is no obvious interpretation of the backbone emission.

Moving type IV bursts involve electrons of higher energy than those produced in bulk energisation. There are at least three different classes of moving type IV bursts (Smerd and Dulk 1971, Stewart 1985), and only one of these, the *advancing front*, is clearly associated with a shock front. Although it is plausible that the electrons are accelerated at the shock front for this class of moving type IV burst, there is no observational evidence that shock acceleration is involved in the other classes of moving type IV burst. The shock acceleration model illustrated in Fig. 1 is favoured for the advancing fronts.

There are two obvious classes of flare continua (e.g., Robinson 1985). The *early flare continuum* (FCE) appears during the impulsive phase, and the *type II-related* (FCII) flare continua are related to associated type II bursts. FCE is associated with type III bursts and microwave bursts, and may be due to the same general distribution of electrons. FCII is presumably due to electrons accelerated or further energised by the type II shock (e.g., Hudson, Lin and Stewart 1982).

The non-flare-associated radio emission implies that electron acceleration occurs in contexts unrelated to solar flares. There is evidence from type I-type III storms of an association between individual type I and type III bursts (e.g., Kai, Melrose and Suzuki 1985), which suggests that type I bursts are generated by electrons with similar characteristics to those that generate type III bursts. The electrons that generate storm type III bursts can be observed *in situ* in the interplanetary medium. It is found that they have energies that can be considerably lower than those involved in flare associated type III bursts, down to around 2 keV (e.g., Lin 1985). There is no detailed theory for the acceleration of these electrons, but it is widely believed that the acceleration is due to localised reconnection sites in an evolving active region.

Solar Cosmic Rays

Observations of energetic electrons and ions from the sun and theories for their interpretation were reviewed by Ramaty *et al.* (1980). There is an extensive literature on possible acceleration mechanisms and no attempt is made to give a comprehensive summary here. The problems that occur in the interpretation concern the details of the acceleration mechanism, the extent and location of storage of the accelerated particles in the corona, and the spatial diffusion and escape of the particles. Recent articles which refer to earlier work concerning the acceleration mechanism include Steinacker, Dröge and Sclickeiser (1989), Schlickeiser and Steinacker (1989) and Miller and Ramaty (1987), Miller, Guessoum and Ramaty (1990) on the acceleration of relativisitc electrons and ions, and Zaitsev and Stepanov (1985) and Smith and Brecht (1986) on the storage and escape of these particles.

A specific theory that accounts for a notable feature of the observations relates to a Bessel function distribution. The Bessel function distribution gives a good fit to the data on nonrelativistic ions, and the same spectrum is implied by a relatively simple theory, e.g., as reviewed by Forman, Ramaty and Zweibel (1986). The following derivation is a summary of that given by these authors.

Consider acceleration in a volume V. Particles are assumed to be injected into V with a momentum $p = p_0$ and to escape from V with a loss rate ζ_L. Then using (2.1) with (2.2) for $v \gg v_A$, the diffusion equation in momentum space becomes

$$\frac{\partial f}{\partial t} = \frac{1}{p^2}\frac{\partial}{\partial p}\left(\zeta_A \frac{cp^2}{4v}\frac{\partial f}{\partial p}\right) + \frac{Q\delta(p-p_0)}{4\pi p_0^2} - \zeta_L f, \tag{5.1}$$

where Q is the number of particles injected per unit time and per unit volume. Steady state solutions of (5.1) may be obtained in the nonrelativistic and in the ultrarelativistic limits. The Bessel function solution is the steady state solution in the nonrelativistic limit. It is

$$f = \frac{2Q}{\pi p p_0 mc\zeta_A}\begin{cases} I_2(z)K_2(z_0), & \text{for } p < p_0, \\ I_2(z_0)K_2(z), & \text{for } p > p_0, \end{cases} \tag{5.2}$$

with

$$z = 4\left(\frac{p}{mc}\frac{\zeta_L}{\zeta_A}\right)^{1/2}, \tag{5.3}$$

with z_0 given by replacing p by p_0 in (5.3). The analogous solution in the ultrarelativistic limit is

$$f = \frac{Q}{4\pi p_0^2\zeta_A(9/16+\zeta_L/\zeta_A)^{1/2}}\left(\frac{p}{p_0}\right)^{-3/2\pm2(9/16+\zeta_L/\zeta_A)^{1/2}} \tag{5.4}$$

with the + sign applying for $p < p_0$ and the − sign applying for $p > p_0$.

The agreement between theory and observation lends support for the validity of a model in which the acceleration is stochastic and reaches a balance between injection and escape. Note that the form of the spectrum for $p \gg p_0$ is insensitive to the form of the injection spectrum at $p \approx p_0$, but that the shape of the escaping spectrum is sensitive to any assumed momentum dependence of escape rate.

6. Creation of Suprathermal Particles from a Thermal Distribution

Most acceleration mechanisms for ions are effective only in accelerating a seed population of already suprathermal particles, and a separate mechanism is required to create this suprathermal population from the thermal population. It is argued here that most mechanisms strongly favour one species over another. As discussed in detail in §7, although the abundances of energetic ions show anomalous features, the most notable characteristic feature is that the relative abundances of suprathermal ions are much closer to those of thermal ions than most acceleration mechanisms imply. This leads to the question of possible mechanisms that can produce suprathermal ions without a strong preference for one species over another.

Formation of a Suprathermal Tail

The formation of a suprathermal tail on a thermal distribution of particles subjected to an acceleration process may be attributed to a balance between the acceleration process tending to produce a nonthermal distribution and collisions (or a collision-like process) tending to maintain the Maxwellian distribution. Both the acceleration and the thermalisation processes may be described by a diffusion equation in momentum space of the form (2.1). Solutions of the appropriate equation were presented by Gurevich (1960), and also by Kruskal and Bernstein (1964) and Knoepfel and Spong (1979) in connection with runaway electrons. An important step is to separate into two regimes, separated by a critical momentum p_c say. In the low momentum regime $p \ll p_c$ the thermalisation process dominates, and in the high momentum regime $p \gg p_c$ the acceleration process dominates. Following Gurevich (1960), it is helpful to consider the flux of particles in momentum space: for $p \ll p_c$ there is no significant flux and a flux towards higher p becomes significant as p approaches p_c; for $p \gg p_c$ this flux approaches an asymptotic value. Given this asymptotic flux in momentum space, one may use it to estimate the rate per unit time and per unit volume at which particles become suprathermal, cf. (3.37).

These ideas were discussed in connection with the runaway acceleration of electron in solar flares by, e.g., Holman (1985*a,b*), Tsuneta (1985), Kuijpers *et al.* (1981). In connection with the acceleration of ions the results have a profound implication: the rate at which ions become suprathermal is a strong function of the ionic species.

Dependence on Ionic Charge and Mass

The dependence of the rate of production of energetic ions on the charge $q_i = Z_i e$ and mass $m_i = A_i m_p$ of the particle was determined by Melrose (1968*b*). The acceleration rate due to stochastic acceleration $\zeta_{Ai} = A_i \zeta_{Ap}$ (p denotes protons) is proportional to A_i in the nonrelativistic limit, and the collision frequency $\zeta_{0i} = Z_i^2 A_i^{1/2} \zeta_{0p}$ is proportional to $Z_i^2 A_i^{1/2}$. As a result the critical velocity v_{ci}, analogous to (3.34), is proportional to $A_i^{-1/2}$:

$$v_{ci} = \tfrac{1}{2} A_i^{-1/2} V_p (\zeta_{0p}/\zeta_{Ap}). \tag{6.1}$$

The rate of production of suprathermal ions then depends on Z_i and A_i according to, cf. (3.37),

$$y_i = \left(\frac{2}{\pi}\right)^{1/2} Z_i^2 A_i^{1/2} \zeta_p e^{-A_i} \exp\left[-\tfrac{1}{2} Z_i \beta (\tfrac{1}{2}\pi - \arctan 2A_i \beta^{-1})\right], \tag{6.2}$$

with

$$\beta = \left(\frac{\zeta_{0p} V_p}{\zeta_{Ap} c}\right)^{1/2}. \tag{6.3}$$

The important qualitative point is that this rate is strongly dependent on the species, and that such strong dependence, although different in form from (6.3), applies to all mechanisms that involve the populations of suprathermal particles being drawn out the tails of their respective thermal distributions. The relatively minor abundance anomalies observed imply that the seed population for second phase acceleration is not produced by drawing ions out the tail of a thermal distribution in this way.

Initial Energisation through Plasma Jets

In view of the extreme sensitivity of the injection spectrum implied by the foregoing discussion, the most notable feature of the abundances of solar and galactic cosmic rays is that they are close to normal cosmic and solar abundances. This requires that the process by which the ions become suprathermal is basically insensitive to Z_i and A_i, but not entirely independent of these parameters.

The simplest type of mechanism, and arguably the only type of mechanism, that can produce suprathermal ions in a way that is insensitive to Z_i and A_i is through the production of fast mass flows. Let us refer to a supersonic mass flow as a *jet*. Once a jet is formed, it propagates until scattering through collisions or through wave-particle interactions destroys it. The resulting scattered particles form an appropriate seed population for a second step acceleration mechanism.

One way of forming jets is through ejection of plasma at about v_A along the separatrices during magnetic reconnection (e.g., Vasyliunas 1975). This involves acceleration by a perpendicular (to $\boldsymbol{B}$) electric field.

7. Anomalous Abundances

The abundances of solar cosmic rays are often 'anomalous' in the sense that they are not exactly the same as ordinary solar abundances. The most notable features are overabundances of the 'heavy' ions from O to Fe and the overabundance of ^{3}He compared with ^{4}He, e.g., the review by Ramaty *et al.* (1980) of earlier data, and Reames, von Rosenvinge and Lin (1985), Mason *et al.* (1986) and Reames (1988), Reames *et al.* (1988). In the so-called ^{3}He-rich flares, the overabundance can be by a factor 10^3 to 10^4. There is an association of ^{3}He-rich events with certain type III electron events, e.g. Reames, von Rosenvinge and Lin (1985), which suggests that the overabundance of ^{3}He is produced in association with the acceleration of the electrons.

Several different mechanisms have been proposed for the production of the overabundance of ^{3}He. Fisk (1978) suggested that preferential acceleration of ^{3}He could occur due to ion cyclotron waves at a frequency intermediate between the proton and ^{4}He cyclotron frequencies. Other suggestions on ways to produce the anomalous abundances were proposed by Möbius *et al.* (1982), Hayakawa (1983), Mullan (1983) and Kocharov and Kocharov (1984), who reviewed the field in detail. The following discussion is essentially a summary of recent ideas proposed by Winglee (1989).

Winglee (1989) explored an injection model for ions based on the acceleration of plasma upflows from the chromosphere in response to the precipitation of electrons during a flare. The acceleration of the ions is attributed to an ambipolar electric field. This electric field arises due to the ambient electrons heated by the precipitating electrons tending to expand and set up a macroscopic charge separation unless the (unheated) ambient ions are dragged along by the electrons. The ambipolar electric field is directed upwards to retard the expanding electrons and accelerate the required upward flow of ions. The resulting flow speeds of the ions are determined by a balance between the acceleration by the electric field and collisional friction, both of which depend on Z_i and A_i. Hence the resulting flow speed also depends on Z_i and A_i. Winglee (1989) showed that the relative flows of different ion species can lead to ion-ion streaming instabilities, and this results in a tendency for the lighter ions to be decelerated and the heavier ions to be accelerated. He argued that this transfer of energy ultimately leads to the abundance differences observed.

Winglee's idea satisfies the basic requirement suggested above that the injection of ions be through mass flows so that the abundances are approximately normal. The specific mechanism for producing the anomalous abundances seems capable of accounting for the gross features of the anomalies. However, the details of how the ^{3}He anomaly is produced seem somewhat contrived. Also there seems no obvious explanation for the correlation between enhanced ^{3}He/^{4}He events and electron events (Reames, von Rosenvinge and Lin 1985).

8. Conclusions

Ideas on the acceleration of fast particles in the solar corona are continuing to develop in two ways: fundamental principles and specific models. (1) Considerable progress has been made in our understanding of the detailed physics of various acceleration mechanisms, notably stochastic acceleration by MHD turbulence, resonant acceleration and acceleration by shock waves. Acceleration by parallel electric fields and during magnetic reconnection are less well understood. (2) Detailed models for specific phenomena involving acceleration are numerous, but there is virtually no specific case where one can say with confidence that a specific model is well established and likely to remain so. A fundamental difficulty is that there are few direct signatures of the acceleration processes.

Further progress in our understanding of acceleration of fast particles in the solar corona requires (i) further development of the theory of acceleration by parallel electric fields and in magnetic reconnection, which are the leading candidates for bulk acceleration, (ii) further modeling of specific acceleration phenomena to provide alternatives to compare and contrast in accounting for existing data, and (iii) new data relating to acceleration, especially high time and space resolution in radio and X-ray observations of the corona and in situ observations of particles in the solar wind.

An important point that requires further modeling concerns the potential available for acceleration by a parallel electric field. There are two opposing opinions. One is that the current limitation argument precludes a current significantly larger than the value $\approx 10^{12}\,$A inferred from vector magnetic field observations, and then to account for the power released one requires a potential $\approx 10^{10}\,$V. This is the more widely held opinion, but there is little supporting evidence for such a large potential drop. The alternative is that the particle data suggest a potential drop $\approx 10^5\,$V, e.g., Lin and Schwartz (1987), and then the power requires a current $\approx 10^{17}\,$A. Such a model would require many oppositely directed current channels so that the *net* current does not exceed the limiting value $\approx 10^{12}\,$A. There is no evidence for such a current pattern, but this may be due to limitations of the available observational instruments. In the present review only the former model has been discussed, and it is desirable to explore the alternative model in at least as much detail in order to compare and contrast them.

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