INTERRELATION OF SOFT AND HARD X-RAY EMISSIONS DURING SOLAR FLARES. II. SIMULATION MODEL

R. M. Winglee,1,2 G. A. Dulk,1 P. L. Bornmann,3 and J. C. Brown4
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ABSTRACT

Recent observations indicated that quasi-static electric fields may be important in the acceleration of the electrons responsible for hard X-ray emissions during impulsive solar flares. Quasi-static fields are also required to drive and to close the currents carried by the accelerated electrons. Two-dimensional (three velocity) electrostatic particle simulations are presented which incorporate the effects of these fields on the particle dynamics as well as those arising from wave-particle interactions induced by the accelerated particles. In the model, the particle acceleration is initiated by a cross-field current in the corona. Because of the limited cross-field mobility of particles in a collisionless plasma, this current can only be closed in the chromosphere where cross-field transport is aided by collisions. As a result, strong quasi-static electric fields are set up which produce strong downward acceleration of electrons in the primary current region. On adjacent field lines, electrons are accelerated up into the corona to provide a return current, but because of the relatively large cross-section of the return current region, the electrons in this region remain relatively low in energy. The induced perpendicular currents are initially provided by the ions, but, as the chromosphere becomes heated by precipitation of energetic electrons, perpendicular electron currents become important, and these currents in turn modify the particle acceleration in the latter part of the flare. The calculated properties of the soft and hard X-ray and microwave emissions from such a system have the following properties: (1) acceleration by quasi-static electric fields and heating via wave-particle interactions produces electron distributions with a broken-power law, similar to those inferred from hard X-ray spectra; (2) heating of the ambient plasma gives rise to a region of hot plasma which propagates down to the chromosphere at about the ion sound speed; (3) the arrival of this hot plasma region in the chromosphere causes the hard X-ray flux to peak due to modification of the cross-field conductivity so that the passage of soft X-ray fronts and the hard X-ray peak are expected to be correlated; (4) the perpendicular heating of coronal electrons is relatively slow, and this can give rise to the observed delay of the microwave peak relative to the hard X-ray peak; (5) heavy ions are preferentially accelerated across the fields as the cross-field currents form, leading to enhancements of heavy ion abundances in the primary current region; this enhanced heavy ion abundance may account for those inferred from soft X-ray line emissions; (6) the ions in the primary current region are then accelerated upward by the same electric fields accelerating the hard X-ray electrons downward; and (7) this parallel and perpendicular acceleration can give rise to Doppler shifts in soft X-ray line emissions similar to those seen during disk and limb flares.

Subject headings: particle acceleration — radiation mechanisms — Sun: flares — Sun: X-rays

1. INTRODUCTION

In a companion paper (Winglee et al. 1991), simultaneous observations of hard X-rays and soft X-ray line emissions were used to explore the dynamics of the electron and ion acceleration. It was shown that in the five flares examined, the hard X-rays during the rise phase and beginning of the decay phase displayed a broken-power-law spectrum which breaks down (i.e., steepens at high energies above about 100 keV). During the latter part of the decay phase, the spectra evolved into a single-power and/or a broken-power law that breaks up (i.e., flattens at high energies). The breaking-down spectra can be produced by the presence of quasi-static parallel electric fields and wave-particle interactions (Lin & Schwartz 1987), with the break energy a measure of the potential drop between the primary energy release region and the chromosphere. The evolution of the spectrum during the decay phase is attributed to the decay of the driving electric fields and the creation and/or trapping of hot electrons in the flaring flux tube by the mirror force.

Certain features of soft X-ray line emissions were shown to be associated with the above changes in the hard X-ray spectrum, providing further evidence for acceleration via quasi-static fields. In particular for the disk flares examined, the velocity of the blueshifted component increases during the rise phase and the beginning of the decay phase, reaching a maximum at about the time the break in the hard X-ray spectrum (and hence electric field) vanishes. For limb flares, there is little evidence for any long-duration blueshifted component since any upward flowing material is moving orthogonal to the line of sight. Instead, the line profile is dominated by a stationary component which shows nonthermal broadening that peaks near the hard (≥50 keV) X-ray peak (i.e., when the electric field is largest).

These observations as discussed in Winglee et al. (1991) are inconsistent with existing chromospheric evaporation models (e.g., Antonucci 1989, Li, Emslie, & Mariska 1989, and refer-
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1984; Reames, von Rosenvinge, & Lin 1985; Mason et al. 1986; Lin 1987; Mason 1987, and references therein). Several mechanisms (e.g., Fisk 1978; Varvoglis & Papadopoulos 1983; Kocharov & Kocharov 1984, and references therein; Winglee 1989) involving wave-particle interactions, such as ion-sound, ion-cyclotron, and ion-ion steaming interactions, have been proposed, and the relative role of these processes can be evaluated by the simulations.

A description of the model is given in §2. The present work extends preliminary results of Winglee, Dulk, & McKean (1989) to include both the rise and decay phases, and it complements that of McKean, Winglee, & Dulk (1990a, b), which investigates thermal processes during the gradual phase of a flare. In §3 the development of the current system in the corona and chromosphere is discussed, and the relation between the induced current and the driving current is derived. Section 4 describes characteristics of the electron acceleration and heating, including the development of broken-power-law spectra, soft X-ray-emitting fronts, and production of the microwave-generating electron population. In §5 the ion dynamics is described, including the rapid perpendicular acceleration of ions which can lead to rapid nonthermal broadening of soft X-ray line and enhancements in heavy ion abundances, and the slower parallel acceleration of the ions, producing a slow-developing blueshifted component. A summary of results is given in §6.

2. SIMULATION MODEL

Because of the complexity, variety, and nonlinearities of the processes involved, two-dimensional (three velocity) electrostatic particle simulations are utilized. Such simulations allow a self-consistent treatment of the particle dynamics and the associated electric fields, including the acceleration produced by quasi-static electric fields and the heating produced by wave-particle interactions. Electrostatic (as opposed to electromagnetic) simulations are used because they are faster and require less memory, and, more importantly, electromagnetic effects for the strongly magnetized plasmas (i.e., \( \omega_p e < \Omega \)) and nonrelativistic particles as considered here tend to be unimportant (Furukawa, Omura, & Matsumoto 1990; Winglee & Kellogg 1990). Two-dimensional spatial variables are important in order to accurately model the current carried by the precipitating energetic electrons and the return currents from the chromosphere (cf. Winglee et al. 1988b).

The particle trajectories and induced electric fields are solved self-consistently on a two-dimensional grid using the algorithm described in the Appendix. Because full particle dynamics is retained, the grid spacing is limited to the smallest plasma scale length (i.e., a debye length). However, the total system size is much longer than the longest plasma scale length, that is, an ion gyroradius. In the following the system size is \( L_x \times L_y = 1024A \times 128A \), where \( A \) is a grid unit and is on an order of a debye length in the corona near the top of the loop.

Reduced ion masses are also utilized in the simulations which allow the ion dynamics to develop fully over periods that can be easily simulated. This approximation also compresses the length scales. In the following the plasma is assumed to consist of two ion species, representing \( H^+ \) and \( \text{He}^{2+} \) with masses \( m_{\text{He}} = 50m_e \) and \( m_{\text{He}} = 100m_e \), respectively (where \( m_e \) is the electron mass). The latter species represents not only ionized He ions but also heavy ions which are nearly fully stripped (e.g., Ca xix and Fe xxv) since they have approxi-
immediately the same \( q/m \) and hence dynamics. It is these latter ions that produce the soft X-ray line emissions, and, as shown in the following, their dynamics can differ from that of the \( H \) ions. This approximation is not restrictive since the ion mass remains sufficiently large to separate their dynamics from that of the electrons.

With the reduced ion masses, \( \Delta \) is approximately equal to 0.4 of an ambient ion gyroradius; for realistic mass ratios and the same energy ion, \( \Delta \) would be a factor of 6 smaller relative to an ion gyroradius. Assuming an initial coronal temperature of a few hundred eV (i.e., order of \( 10^6 \) K), a magnetic field of a few hundred gauss, and a preflare coronal density of \( 10^9 \) cm\(^{-3} \), the system length when normalized to an ion gyroradius is of the order of several tens of meters by about 10 m. This distance is about \( 10^{-5} \) smaller than a typical flaring loop. The duration of the simulations when normalized to ion time scales is also about \( 10^{-5} \) smaller than a typical flare.

This small system size and duration may appear as a daunting limitation of the work. We argue that it is not for four reasons. First, much of the particle acceleration is determined by processes that occur on plasma scale lengths and not the length of the flux tube. In particular, reconnection is expected to occur over scale lengths of the order of an ion gyroradius (e.g., Priest 1982), as is acceleration via shocks and double layers. The flare currents within the loop are also expected to filament on similar scale lengths (e.g., Winglee et al. 1988b), and these filaments are seen in the auroral zone (e.g., Hoffman, Sugira, & Maynard 1985). The simulations presented are more than sufficient to resolve these processes. Second, the hard X-ray spectra from flares show similar characteristics in common with that from the auroral zone (e.g., Lin & Schwartz 1987), yet the ratio of ion gyroradius to the flux tube length (\( \sim \Delta/L \)) differs by nearly three to four orders of magnitude; this difference in \( \Delta/L \) is similar to the simulation model. Third, while there is compression of the length and time scales, their ordering is maintained. In particular, the duration of the simulations is comparable to a transit time of accelerated (several keV) ions (as in actual flares), and this time scale is much longer than those of most plasma processes. This ordering allows the particles to experience similar types of potentials as in an actual flare while incorporating the effects of wave-particle interactions in a consistent manner. Fourth, testing the predictions against a comprehensive set of observations at multi-wavelengths such as those presented in the companion paper provides stringent testing of any model.

The initial plasma in the system is assumed to be homogeneous across the magnetic field (i.e., in the \( y \) direction). Along the magnetic field (i.e., \( x \)) the density is assumed to increase so as to simulate the change in density from corona to the chromosphere (see Krüger 1979; Vernazza, Avrett, & Loeser 1973) and aid current closure. As shown schematically in Figure 1, the corona and upper transition region are represented by the left portion of the simulation system, with the density increasing approximately by a factor of 4 from near the left-hand boundary (top of loop at \( x/\Delta = 0 \)) to the point \( x = 640\Delta \). The plasma is collisionless until the point \( x = 524\Delta \), where the plasma first becomes weakly collisional, that is near the beginning of the transition region (the implementation of these collisions in the simulations is discussed in the Appendix). The electron density at this point is denoted as \( n_e \) (i.e., the maximum initial density of the corona) and is of the order of \( 10^{10} \) cm\(^{-3} \). The corresponding plasma frequency \( \Omega_{pe} \) is about \( 2\pi \times 900 \) MHz. The assumed magnetic field strength is assumed to be such that \( \Omega_e = 1.5\Omega_{pe} \).

The density then increases by another factor of 5 across the upper chromosphere which becomes increasingly collisional. This collisionality provides enhanced cross-field conductivity, allowing closure of the current system. With this profile, the density increases by a total factor of 20 from the corona to the chromosphere. While this increase is smaller than in the Sun, the results presented here are not sensitive to the assumed density profile provided that the chromospheric density is much greater than the coronal density; simulations where the density ratio is as high as 100 and as small as 10 show essentially the same results.

From 950\( \Delta \) to 1020\( \Delta \), the density is ramped down so that the plasma is not actually in contact with the right-hand boundary. This decrease reduces the influence of any boundary conditions on the simulations, although this proved to be an unnecessary precaution. Also the chromospheric plasma density is known to decrease and the neutral density to increase as the photosphere is approached. A total of about 250,000 particles are used in the simulations. To save computer
memory, particles at the lower altitudes have a higher charge and mass (but the same charge-to-mass ratio), which allows the high-density regions to be represented by about the same number of superparticles as the tenuous plasma at higher altitudes, without changing the particle dynamics.

The thermal velocity of the coronal plasma $v_{T\rho}$ at the highest altitude (i.e., near the left-hand boundary) is assumed to be equivalent to about 100 eV. The temperature across the rest of the system, except in the region of decreasing density in the chromosphere, is then varied so that the plasma is initially in pressure equilibrium, that is, $nT$ is constant along the system except near the right boundary. In this latter region, the temperature is assumed to be constant, being equal to the minimum chromospheric temperature. Pressure balance here is provided by the presence of enhanced collisions arising from increased neutrals. Throughout the system the relative concentration of the He is kept constant at 12.5% of the total ion density.

3. THE DEVELOPMENT OF THE CURRENT SYSTEM

The driver for the current system and associated particle acceleration assumed here is a cross-field current at high altitudes in the corona. The origin of this current and how it affects the coronal and chromospheric plasma is shown schematically in Figure 2. The driving current is envisaged to arise from current sheet acceleration in the primary energy release site (Fig. 2a). Such current sheets which are supported by an electric field directed perpendicular to the magnetic field are central to flare models driven by reconnection (e.g., Priest 1982). This electric field is expected to accelerate preferentially the ions (compared with the electrons) perpendicular to the magnetic field, particularly if there is a weak magnetic field normal to the current sheet (e.g., Lyons & Speiser 1982). This same electric field causes the plasma outside the current sheet to convect inward, so that there is continual inflow and acceleration of ions. It is proposed here that after this acceleration occurs and the current-sheet electric field decays (so there is no force confining the plasma to the current sheet), these energetic ions will tend to flow preferentially outward across the ambient magnetic field since their energy and gyroradius are larger than those of the electrons (Fig. 2b). It is this flow which produces the driving current assumed in the present model and which in turn produces the field-aligned potential (Fig. 2c) and subsequent acceleration of coronal and chromospheric plasma. This mechanism may also be responsible for the formation of the current wedge seen at substorm onset (McPherron, Russell, & Aubry 1973); it is addressed in more detail in another paper (Winglee 1991).

This cross-field current is incorporated into the simulation system by injecting both electrons and energetic ions from the central portion of the system near the left-hand boundary (the actual processes associated with the current sheet acceleration and associated cross-field current are not included explicitly although a simulation model with these processes included is currently under development). The injection occurs over a width of 32Å across the magnetic field, which is equivalent to about 14 gyroradii of an ambient H ion or about one gyroradius of a 20 keV H ion. This width is about the maximum size that the plasma can support without current filamentation occurring (Winglee et al. 1988b). The electrons with their small gyroradius remain tied to the field lines and carry little perpendicular current (Fig. 3a). The energetic ions, on the other hand, have a much larger gyroradius and stream across the field lines to set up the cross-field current, similar to the above description of the decay of the current sheet. Thus, in the center of the system (hereafter called the primary current region), there is an excess of electrons, while on adjacent field lines (hereafter called the return current region), there is an excess of positively charged ions.

The electrons with their relatively light mass are the first to respond to the applied current as shown schematically in Figure 3a. Electrons in the primary current region are accelerated down into the chromosphere while those in the return current region are accelerated upward into the corona. This
Fig. 3.—Schematic of the driving current in the primary energy release region and the induced electron and ion currents. The electrons are the first to respond to the driving current, and strong field-aligned currents are induced; little perpendicular current is carried by the electrons until they reach the chromosphere where collisions enhance their cross-field mobility. The ions, with their relatively large mass and gyroradius, respond later to the driving current and are initially accelerated across the field lines, providing some of the parallel current needed for current closure (Fig. 3b). These ions, once in the beam region, are accelerated up into the corona by the same electric field accelerating the electrons downward.

The energies to which the particles are accelerated depends on the magnitude of the driving current (the actual dependence is derived shortly). This dependence allows the rise and decay phases to be modeled by applying a triangular time profile for the driving current. In the following, the maximum value of the driving current flowing from one side of the beam region into the return current region has been chosen to be of the order of the total integrated "thermal ion current" along the flux tube, that is \( \sim n_e v_{Th} L_x \) where \( v_{Th} \) is the thermal velocity of the coronal H ions. This current cannot actually be supported by the ambient plasma without significant acceleration being required since gyromotion limits the net flux of coronal ions across the field lines. In particular, at the earlier times, \( \omega_{pe} t \lesssim 400 \) (Fig. 4a), the induced perpendicular currents are primarily ion currents and amount to less than half the driving current.

Because of this lag in the balancing perpendicular current, parallel currents develop along the field lines in an effort to alleviate the charge imbalance, as seen in Figures 4b and 4c. These parallel currents are of the order of a few electron thermal currents (i.e., \( n_e v_{Th} \)) in the primary current region and about one-third smaller in the return current region; they are present through the bulk of the system, with strong currents being driven in the chromosphere at \( x / \Delta \approx 600 \) at \( \omega_{pe} t \gtrsim 400 \) (even though the actual driving current is confined to near the left-hand boundary; see Winglee et al. 1988b).

These induced currents can be related to the driving perpendicular current and the assumed geometry by current continuity as follows. The total driving perpendicular current in Figure 4a is of the form

\[
I_D = 2a(t) n_e v_{Th} L_x ,
\]

where \( a(t) \) is a time-dependent measure of the magnitude of the driving current and \( n_e \) is the maximum initial density of the coronal electrons. (The factor of 2 arises since there are outward currents from both sides of the beam region.) Since the primary current region provides one of the current elements between the driving current and the chromosphere, then, assuming current closure, the integrated parallel current across the beam or primary current region \( I_{\|b} \) is equal to \( I_D \), that is

\[
I_{\|b} = n_e v_{Th} L_c \approx I_D ,
\]

where \( v_b \) is the "bulk velocity" of the current carriers in the beam region and \( L_c \) the width across the field lines of the beam region. (The actual bulk velocity can be different locally since the current density in equation (2) is normalized to the parameter \( n_e \) and not the local density.) Similarly, since the return current region provides the matching current element to the beam region, current balance requires that

\[
I_{\|r} = n_e v_c L_r \approx I_D ,
\]
where the subscript $r$ refers to quantities of the return current region.

On substitution of equation (1) into equations (2) and (3), the “bulk velocities” in the beam and return current regions are given by

$$v_b \simeq 2\alpha(t) \left( \frac{m_e}{m_H} \right)^{1/2} \frac{L_x}{L_b} v_{T_e} \tag{4}$$

$$v_r \simeq 2\alpha(t) \left( \frac{m_e}{m_H} \right)^{1/2} \frac{L_x}{L_r} v_{T_e} \tag{5}$$

In the present example, $L_x/L_b = 32$, $L_x/L_r = 10.7$, and $(m_e/m_H)^{1/2} = 7$ so that $v_b \simeq 4.5v_{T_e}$ and $v_r \simeq 1.5v_{T_e}$ when the driving current reaches its maximum. These estimates give current densities similar to the numerical results shown in Figure 4c.

Thus, from equations (4) and (5), the strongest particle acceleration in the primary current region occurs when the current elements are long and thin, that is, $L_x/L_b \gg 1$. A similar criteria is also true for the return current region so that if the return current regions are relatively wide (i.e., $L_x/L_r \ll 1$), the return current electrons are not energetically important. Only in cases where the filling factor is of order unity are the return current electrons moving up the field lines with energies similar to those of the electrons in the primary current region.

One other important characteristic of the current system is that at times $\omega_{pe} t \gtrsim 400$, the electron perpendicular current becomes dominant (Fig. 4a). At the same time, parallel currents are seen to penetrate deep into the chromosphere (i.e., at $x/\Delta \gtrsim 900$). As shown in the following section, the change in current characteristics is due to the precipitation of energetic electrons into the chromosphere and subsequent modification of the chromospheric conductivity.
4. ELECTRON DYNAMICS

As the above current system develops, the electrons are accelerated and heated by both the quasi-static electric fields driving the currents and by wave-particle interactions induced by the currents. The characteristics of the heating and acceleration in the primary and return current regions are presented in this section. Because of the different sizes of the primary and return current regions, the heating and acceleration in the two regions are quantitatively different (§ 4.1), and this nonuniformity has important consequences for the processes responsible for soft X-ray emissions (§ 4.2) and the microwave and hard X-rays emissions during solar flares (§ 4.3).

4.1. Electron Acceleration and Heating

The downward acceleration and heating of electrons in the primary current region are illustrated in Figure 5 which shows the evolution of the parallel velocity $v_{pe}$ of the beam electrons (i.e., those electrons injected as part of the driving current) as a function of distance along the magnetic field. The left-hand side shows the development during the rise phase, and the right-hand side, the decay phase.

The initial energies of the beam electrons are similar to those of the ambient electrons (i.e., $v < 3v_{Te}$). Through their interaction with the ambient electric field, these electrons are accelerated to higher and higher velocities during the rise phase and the beginning of the decay phase. In particular, the distribution of the electrons near the left-hand boundary at $x/A \approx 300$ is seen to change from one which has no net drift to one with a net streaming velocity approximately equal to $2v_{Te}$ in Figure 5b. This streaming velocity increases as the driving current is increased, reaching a maximum value of about $17v_{Te}$ near the peak of the driving current (Figs. 5e and 5f). During the decay phase, the streaming velocity decreases with decreasing driving current.

This streaming velocity can be equated to the “bulk velocity” defined by equation (4) by noting that the bulk velocity is defined as the average electron velocity required to produce the prescribed current density for a plasma of density $n_e$. Since the actual coronal density at high altitudes $n_{ex}$ is smaller, the streaming velocity $v_s$ must be larger by a factor of $n_e/n_{ex}$, that is

$$v_s = n_e/n_{ex} v_b = 2ax(t) \left( \frac{n_e}{n_{ex}} \right) \left( \frac{m_e}{m_H} \right)^{1/2} \frac{I_p}{L_x} v_{Te}. \quad (6)$$

At high altitudes in the simulation, $n_e/n_{ex} \approx 4$ so that at the current peak where $v_b \approx 4v_{Te}$, equation (6) predicts that $v_s \approx 16v_{Te}$, similar to that seen in Figure 5. Equation (6) is approximately satisfied during both the rise and decay phases. Note that assuming an initial coronal temperature of 100–200 eV (i.e., $1-2 \times 10^6$ K), the above streaming velocity corresponds to an energy of the order of 20–40 keV, that is, to hard X-ray-producing energies.

In addition to the bulk acceleration, there are waves induced in the plasma as the beam propagates through the plasma. Their spectrum is similar to that in Winglee et al. (1988); it is not shown but includes intense ion cyclotron, lower hybrid, and Langmuir waves. The presence of the Langmuir waves (which are produced by a beam interaction between the downflowing energetic electrons and the ambient plasma) can be seen in Figures 5c–5f, as vortices in phase space. As these waves grow, some electrons are scattered to low energies, while others are scattered to velocities nearly twice the streaming velocity. In other words, through the wave-particle interactions, a population of electrons is produced with energies nearly 4 times higher (in the present case to energies of 80–180 keV).

The phase spaces in Figure 5 show two other important features. First, as the Langmuir waves propagate into the collisional plasma, the low-energy electrons with their high collision frequency are subject to more drag than the high-energy electrons, causing the breaking of the coherent vortices (e.g., Figs. 5d and 5e). Second, the loss of energetic electrons during the decay is relatively fast, with few energetic electrons being present after the driving current has been turned off (Fig. 5j). This rapid loss is due to the fact that the magnetic field in the present model is assumed to be homogeneous so that there is no mirror force to confine even intermediate energy electrons in the corona. Instead, they precipitate almost directly into the chromosphere where their energy is dissipated.

The response of the ambient electrons as the current system develops during the rise phase is shown in Figure 6. The decay phase (not shown) is essentially the reverse of that in Figure 6 except that the final temperature of the plasma is about 2–4 times hotter than the original plasma. It is seen that strong electron acceleration and heating occurs rapidly along the length of the system. For example, in Figure 6b there are electrons being accelerated into the transition region at $x \approx 600 \Delta$.
with speeds nearly twice the maximum speed of the initial plasma in the region. These electrons appear in this region well before any of the beam electrons (e.g., Fig. 5b). This acceleration is due to the rapid propagation of the electric field supporting the current system. One consequence is that electrons in the low corona, while experiencing only a small fraction of the potential drop between the primary energy release region and the chromosphere, only have a small distance to travel to the chromosphere and are thereby able to arrive there earlier than the more energetic electrons which originate near the primary energy release region at the top of the loop. In other words, low-energy phenomena such as Hα emissions are predicted to start before more energetic emissions such as the hard X-rays. This prediction differs from a thermal model where the energetic electrons are produced in a local region with a thermal distribution; the faster, more energetic electrons in this case essentially free-stream down the field lines and reach the chromosphere first.

The evolution of the total velocity distribution for electrons in the primary current region is shown in Figure 7. Distributions in the upper, middle, and lower corona are shown across the figure at 5 times during the flare. The distributions are taken from a region δx = ±50A from the central position, with consecutive contours differing by a factor of 2. In the upper corona, a well-defined beam feature at v_x ≈ 12v_t is seen to develop quickly (Fig. 7b). There is also a stationary component but one which is relatively tenuous, as seen from the contour lines. At early times (Figs. 7a and 7b), most of the electrons experience strong acceleration in the parallel direction but gain little perpendicular energy. This lack of perpendicular acceleration arises because the small gyroradius of the electrons inhibits acceleration across the field lines. Instead, they are subject to an $E \times B$ in the $z$ direction. Energetically, these drifts are unimportant, but, through wave-particle interactions (especially from lower and upper hybrid waves), some of these particles are eventually scattered to high perpendicular energies (Fig. 7c).

It is seen from the distributions that this perpendicular heating is most efficient for electrons with moderate parallel velocities, rather than for the fastest electrons. The largest perpendicular velocities (about 100v_t) are attained at about $\omega_{pe} t = 1080$, which is after the peak of the driving current. This delay is due to the fact that free energy of the wave-particle interactions is still being provided by the induced quasi-static field (even though the absolute value of the fields is decreasing). This delay, as discussed in § 4.3, may be important in explaining the observed delay between the microwaves and the hard X-rays.

An interesting feature of the distributions in the upper corona (Figs. 7b–7e) is that there is actually a hole (depletion) at $v_x \approx v_\perp \approx 3v_t$. This hole feature arises from pitch angle scattering associated with the wave-particle interactions producing the perpendicular heating. The resulting distribution is unstable to the electron-cyclotron maser instability (Wu & Lee 1979) which is believed responsible for microwave spike bursts (Melrose & Dulk 1982). Previously, loss-cone distributions produced by mirroring and precipitation in convergent field lines have been invoked to drive the instability. The present work indicates that such loss-cone distributions are not necessary; instead the maser emission can be driven unstable by distributions resulting from the imposition of $E \times B$ drifts associated with localized current systems (see Wingee & Kellogg 1990). These emissions should be strongest slightly after the hard X-ray peak where the electron distribution shows large positive $v_x$ gradients.

Further down the field lines where the initial plasma density is a factor of 2 higher (middle panels), the high-energy electrons first appear as a high-energy tail (Fig. 7g). Eventually a beam feature develops (Fig. 7h) with its density comparable to that of the stationary electron component. Again the presence of strong perpendicular heating appears relatively late and is long lasting (Figs. 7h–7j). Near the transition region and chromosphere where the ambient density is relatively high, the high-energy electrons appear only as a high-energy tail.
current regions are shown in Figure 8. The distribution in the upper and middle corona show bulk acceleration to speeds only a few times the ambient thermal speed; high-energy electrons appear primarily as tails, and there is little evidence of perpendicular heating. In the lower corona, there is less bulk acceleration since the plasma density is relatively high and the return current electrons appear as a tail in the distribution. In other words, throughout most of the return current region, the electrons remain relatively cool compared with the electrons in the primary current region, consistent with the analysis in § 3.

4.2. Implications for Soft X-Ray Emissions

From the distributions in Figures 7 and 8, it is clear that the electron distribution cannot be accurately described by a simple temperature, as is sometimes assumed in interpreting X-ray emissions. At each spatial location, the electron distribution can be represented by a stationary component and a high-energy component which may or may not be of higher density than the stationary component. The high-energy component can lead to ionization of heavy ions, such as Ca xix and Fe xxv, whose soft X-ray line profiles are used to infer the outflow of chromospheric plasma. On the other hand, the low-energy stationary component can produce recombination with any ionized ion.

Thus, as a first approximation to the evolution of both the production of highly ionized heavy ions and the soft ($\gtrsim 10^7$ K) X-ray-emitting plasma, the effective electron density $n_x$ involved in these processes can be approximated by the density of electrons with energies above 1 keV minus one-third of the density of electrons below 1 keV. Regions where this density is positive have electrons of sufficient density and energy that Ca xix is the dominant ionization state and the density gives a measure of the soft X-ray intensity. More detailed calculations
using a seven-component electron distribution to represent those in Figures 7 and 8 show that this simple subtraction method gives a good approximation to regions where Ca xix is the dominant ionization state.

The evolution of this hot plasma region is shown in Figure 9 (only regions of positive effective density are included). It is seen that through the bulk of the flare, Ca xix is the predicted dominant ionization state of Ca in the primary current region. Ca xix is also produced in the return current regions but mainly during the decay phase and at relatively low density. Thus, the soft X-ray emission is expected to be dominated by processes in the primary current region. During the rise phase, the hot plasma "propagates" from the primary energy release region down to the chromosphere where it stops. The word "propagation" is used here because the electrons that comprise the hot plasma are not constant in time. Indeed, energetic (>10 keV) electrons are continually flowing through the region, impacting on the chromosphere well before this hot region moves down into the lower corona.

The region of highest density, and hence soft X-ray emissions associated with this hot plasma region, is seen to propagate downward through the corona during the rise phase (dotted line) with a speed of the order of 1.5\(v_T\). For the assumed ratio of \(m_e/m_i\) in the simulations, this speed corresponds to an ion sound speed of a plasma at about \(2 \times 10^7\) K, that is, the average energy of the hot electrons. Since this is the region where the soft X-ray intensity is largest, soft X-ray imaging would show this material as a downward propagating front with a speed comparable to the ion sound speed of a hot plasma at a few million degrees.

This downward propagation of the soft X-ray-emitting material is similar to that observed by Nitta, Kiplinger, & Kai (1989). Soft X-ray-emitting fronts have also been reported by Rust, Simnett, & Smith (1985), and Batchelor et al. (1985) have reported a correlation between the hard X-ray rise time and the propagation time of fronts moving at the ion sound speed. The above results are consistent with these observations.

During the decay phase, the position of the hot plasma is seen to retract back up to the upper corona after the hard X-ray emissions have ceased (i.e., after about \(\omega_{pe}t = 1560\)), similar to the observations of Nitta et al. (1989). Once this retraction begins, its apparent speed is comparable to the expansion speed during the rise phase. This speed is determined by the loss rate of electrons to the chromosphere. Since there are no convergent magnetic fields in the present model to confine the hot plasma in the corona and limit losses to the chromosphere, this speed is an overestimate of that likely to occur during the decay phase of a flare.
Fig. 9.—Contours of the density of the hot region where the bulk of the electrons have energies above 1 keV during the rise phase (left-hand side) and decay phase (right-hand side). This region, which is important in determining the predicted soft X-ray emissions, is primarily confined to the primary current region and propagates down into the chromosphere during the rise phase at a speed of the order of the ion sound speed.

4.3. Implications for Hard X-Rays and Microwaves

Another test of whether the model is accurately describing the dominant processes involved in flare acceleration is in the predicted energy spectra of the electrons precipitating into the chromosphere. Figure 10 shows the computed energy spectrum ($E > 100$ eV) of electrons below the top of the transition region with energies at $x/\Delta > 524$. The few electrons that appear initially primarily reside in the upper transition region. As electrons are accelerated and heated by the electric fields associated with the driving current, additional low-energy ($100$ eV–$1$ keV) electrons first start precipitating into the transition region (Fig. 10a). These electrons, as discussed in § 4.1, were initially lying in the low corona; they can produce low-energy phenomena such as Hz emissions before the rise of the hard X-rays in the low corona. It is for this reason that this phase is indicated as early rise phase. A statistical study of 93 hard X-ray flares observed by the Solar Maximum Mission (SMM) shows that the early arrival of low-energy electrons in the chromosphere is a common occurrence (Dulk et al. 1991).

At later times during the rise phase, the spectrum in the X-ray energy band is seen to extend to higher energies and to harden. This hardening is due to increasing acceleration and heating of a higher percentage of electrons along the flaring flux tube, as seen in the phase spaces in Figures 5 and 6.
Fig. 10.—Spectrum of the energetic electrons in the transition region and chromosphere, i.e., at $x/A > 524$. The first energetic electrons to appear have energies of only a few hundred eV and initially reside low in the corona. At later times the spectrum hardens and at hard X-ray energies has a broken-power law distribution (e.g., dashed line).

Moreover, during the rise phase, the spectrum over hard X-ray energies displays the characteristics of a broken-power law (dotted lines), with the break energy corresponding to the energy associated with the streaming velocity given by equation (6) (or equivalently the potential of the electric field in the primary current region), similar to the suggestions of Lin & Schwartz (1987) and Winglee et al. (1991). During the decay phase, it develops a single-power law over hard X-ray energies. A breaking-up spectrum is not seen as in the observations. This discrepancy in the decay phase is attributed to the rapid loss of electrons into the chromosphere because of the assumption of straight field lines, as discussed in § 4.2.

Another important feature predicted by the simulation model is the overall profile of the hard X-rays in relation to both the propagation of the soft X-ray-emitting plasma and microwave emissions. This timing can be inferred from the time histories of the parallel and perpendicular energies ($E_\parallel$ and $E_\perp$) of the beam, coronal, and chromospheric electrons shown in Figure 11. Since most of the high-energy electrons are propagating downward, $E_\parallel$ of the beam and coronal electrons approximately represents the energy available for the production of hard X-rays. The energy is seen to peak at about $\omega_{pe} t \approx 620$, which corresponds to the time when the soft X-ray-emitting plasma reaches the chromosphere (§ 4.2) and is before the maximum of the driving current. This association between the soft and hard X-rays is due to the fact that the...
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The parallel energy of the beam and coronal electrons (which is an indicator of the energy available for hard X-ray production) rises rapidly, reaching its peak at about $\omega_{pe}t \approx 840$. The perpendicular energy rises more slowly, reaching its peak at about a few tens of seconds, then the above delay is of order of a few seconds, which would be comparable to the observed delay.

5. ION DYNAMICS

The electric fields required to produce current closure cause the acceleration of the ions to energies of several keV on short time scales, with preferential acceleration of heavy ions into the primary current region (§ 5.1). As these ions enter the primary current region, they are accelerated upward by the same electric field that accelerates the primary electrons downward (§ 5.2). The characteristics of the acceleration are shown (§ 5.3) to be consistent with those inferred from soft X-ray line emissions, including the line broadening and Doppler shifts relative to the hard X-ray emissions as described by Winglee et al. (1991).

5.1. Perpendicular Ion Acceleration and Heavy Ion Abundances

The ions are subject to fast, perpendicular acceleration. This is illustrated in Figure 12 which shows the $v_y$ component of the $H^+$ and $He^{++}$ ions as a function of position along the magnetic field during the first half of the rise phase. It is seen that there is a region of enhanced perpendicular ion energies which "propagates" from $x/A \approx 200$ at $\omega_{pe}t = 120$ to $x/A \approx 500$ at $\omega_{pe}t = 240$ and then into the chromosphere by $\omega_{pe}t = 360$; further "propagation" is slowed by collisions in the chromosphere. This "propagation" corresponds to the motion of the beam (or hard X-ray-producing electrons) through the plasma (viz., Fig. 5) and the development of the current system along the field lines; the ions themselves do not propagate at this fast speed.

This perpendicular acceleration can arise from perpendicular acceleration (in $y$) and subsequent gyromotion and/or from induced $E \times B$ drifts. If it is due to $y$ acceleration, then different ion species should have similar energy, that is, $v_y^2/v_x^2$ should be the same for the different ion species (where $v_y$ is the thermal speed of a given ion species). On the other hand, if it is due to $E \times B$ drifts, then the maximum $v_y$ should be the same in absolute terms. In Figure 12 it is seen that the $H$ ions reach maximum velocities of only about $4v_{Th}$ (except close to the primary energy release region), whereas the $He^{++}$ ions reach velocities as high as about $7v_{Th}$. In absolute terms, the maximum velocities of the ions are approximately the same (noting that $v_{Th} = \sqrt{2kT_m}$). Thus, the primary source of the ion perpendicular energy is from induced $E \times B$ drifts.

The magnitude of the $E \times B$ drift velocity $v_p$ can be equated

arrival of the hot plasma at the chromosphere enhances the cross-field conductivity and thereby reduces the need for large electric fields to close the current system. This reduction in electric field in turn limits the acceleration of hard X-ray-emitting electrons.

Note also that the perpendicular energy of the beam and coronal electrons (which is a measure of the energy available for microwave production) rises more slowly than the parallel energy and is also slower to decay. This lag, as discussed in § 4.1, is due to the fact that the perpendicular heating is being driven by wave-particle interactions rather than direct acceleration and that these interactions can continue provided the driving current is nonzero. Further, the perpendicularly heated electrons are lost much more slowly from the flux tube than the hard X-ray-emitting electrons since they tend to have smaller parallel velocity. The addition of convergent field lines would further reduce their loss rate.

This difference in the characteristics of parallel and perpendicular energies may have important consequences for microwave emissions in relation to the overall flare dynamics. In particular, at early times the energetic electrons are primarily directed downward along the field lines so that little microwave emission is expected. Because the energetic electrons (Fig. 7) have such small pitch angles, this result is not expected to change for convergent magnetic fields with typical mirror ratios of $B_{max}/B_{min} \leq 6$. Strong microwave emissions are only expected in the present model after the electrons have experienced strong perpendicular heating which, as discussed earlier, does not begin to occur until late in the rise phase and beginning of the decay phase. Indeed, from Figure 11, it is seen that the electron perpendicular energy does not peak until about $\omega_{pe}t \approx 840$, which is after the peak of the parallel energy and hence hard X-rays.

This delay is qualitatively in agreement with the delay between the microwave hard X-ray spikes observed by Cranell et al. (1978), Kaufmann et al. (1983), and Cornell et al. (1984). Moreover, if the above time scales are normalized to the rise time of a flare (i.e., a few tens of seconds), then the above delay is of order of a few seconds, which would be comparable to the observed delay.

Fig. 11.—Time histories of the parallel and perpendicular energies of the beam, coronal, and chromospheric electrons, normalized to the initial perpendicular energy of the coronal electrons $E_0$. The parallel energy of the beam and coronal electrons (which is an indicator of the energy available for hard X-ray production) rises rapidly, reaching its peak at about $\omega_{pe}t \approx 620$. The perpendicular energy (which is an indicator of the energy available for microwave production) rises more slowly, reaching its peak at about $\omega_{pe}t \approx 840$. Strong microwave emissions are only expected in the present model after the electrons have experienced strong perpendicular heating which, as discussed earlier, does not begin to occur until late in the rise phase and beginning of the decay phase. Indeed, from Figure 11, it is seen that the electron perpendicular energy does not peak until about $\omega_{pe}t \approx 840$, which is after the peak of the parallel energy and hence hard X-rays.

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The magnitude of the $E \times B$ drift velocity $v_p$ can be equated
Fig. 12.—Evolution of the orthogonal velocity component $v_x$ of the H ions (left-hand side) and the He ion (right-hand side). Increases in this velocity component (which can arise from $E \times B$ drifts and/or perpendicular acceleration and gyration) are seen to develop rapidly and propagate down the field lines at speeds much faster than that of the hot plasma region in Fig. 9. This rapid development is due to the propagation of the current system rather than the particles themselves.

to the characteristics of driving current as follows. By definition,

$$v_D = \frac{E \times B}{B^2} c.$$ 

The electric field strength in equation (7) can be approximated by

$$E \approx \phi/(L_b) \approx 0.5m_e v_x^2/L_b,$$ 

where it has been assumed that (1) the potential drop along the magnetic field is approximately equal to the potential drop across the field lines between the primary and return current regions, and (2) this potential drop occurs over a distance of the order of the width of the primary current region $L_b$. On substitution of equation (6) into equation (8), the expression for $v_D$ reduces to

$$\frac{v_D}{v_{TH}} \approx 2\omega t \left( \frac{m_e}{m_H} \frac{n_e}{n_H} \right)^{1/2} \left( \frac{L_x}{L_b} \right)^{1/2} \frac{\rho_H}{L_b},$$

where $\rho_H$ is the H gyroradius. For the parameters considered here, equation (9) predicts $v_D \approx 3v_{TH}$, which is consistent with that in Figure 12.

While there are large $E \times B$ drifts induced in the ion motion, there are also strong $y$ accelerations, particularly for the heavy ions, and as well as thermalization of the $E \times B$ drifts by wave-particle interactions. These processes are illustrated in Figures 13 and 14 which show the $v_y - y$ and $v_z - y$ spaces for the He ions in the corona during the rise and decay phases, respectively. The induced $E \times B$ drift can be seen in $v_x$, as the bulk acceleration of the ions to large positive $v_x$ for ions with $y/\Delta > 64$ and large negative velocities for $y/\Delta < 64$. These bulk drifts are maintained through the rise phase and the beginning of the decay phase.

The $v_y$ component shows different characteristics. During the rise phase, there is net ion flow into the primary current region, ions with $y > 64\Delta$ having negative bulk velocities and those ions with $y < 64\Delta$ having positive velocities (e.g., Figs. 13b and 13d). This net flow is also seen as an increase in the phase space density in the middle of the system. During the decay phase, the electric field (which initially accelerated the ions inward) decreases, and many of the energetic ions which were initially accelerated into the primary current region are

Fig. 13.—The perpendicular $v_y$ and orthogonal $v_z$ velocity components of the He ions in the region $256 < x/\Delta < 512$ as functions of $y$ during the rise phase. Their motion is seen to comprise a net inward drift in $y$ (i.e., $v_y < 0$ for $y/\Delta > 64$ and $v_y > 0$ for $y/\Delta < 64$), oscillations in $v_z$, and strong $E \times B$ drifts in $v_x$. 

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now able to diffuse back out into the return current regions (Figs. 14b and 14c).

Superposed on this net flow are oscillations. For example, in Figure 13c the $y$ flows are in the opposite direction to those in Figures 13b and 13d. The period of these oscillations is approximately the He cyclotron frequency. The evolution of the $v_y - y$ phase space of the H ions (not shown) displays a similar behavior, except that the dominant oscillation frequency is the lower hybrid frequency. These oscillations are due to the sloshing of the ions about the magnetic field under the influence of the perpendicular quasi-static electric field. They have three effects. First, due to the higher oscillation frequency of the H ions, their net flow into the primary current region is smaller than that of the heavy ions. Second, the associated wave-particle interactions scatter some of the ions, producing high-energy tails. These tails are particularly evident during the decay phase as seen in Figure 14. Third, these oscillations lead to modification of the ion charge density in the primary current region, which in turn modifies the structure of the electric potential down the field lines and the associated electron acceleration. This modulation of the electron acceleration is seen in the time history of the perpendicular electron current in Figure 4 and in the changes in the energy of the beam electrons near the hard X-ray peak (Figs. 5e and 5f).

The preferential flow of heavy ions into the primary current region is illustrated in Figure 15, which shows cuts of the ion charge density profile in the center of the primary current region (top panels) and return current region (lower panels) at $\delta y/\Delta = 40$ from the center of the system. While there are some oscillations in the density associated with the above wave-particle interactions, the H density down to the middle corona is seen to increase by about 20%. This increase is accompanied by...
by a decrease in density in the return current region as the ions flow across the field lines. At lower altitudes, larger increases of the density of the primary current region are seen because of outflow of chromospheric plasma. The density of the He ions shows similar increases in the primary current region, except that enhancements of nearly 100% are typical; that is, the relative density of heavy ions in the primary current region is enhanced. During the decay phase of the flare (not shown) there is a tendency for the ions to diffuse out of the primary current region, as discussed above.

One important aspect of this preferential acceleration of heavy ions into the primary current region is its effect on heavy ion abundances. Sylwester, Lemen, & Mewe (1984) and Sylwester (1987), using the characteristics of the soft X-ray line profiles, showed that the abundance of heavy ions increases during the rise phase and decreases during the gradual phase. This apparent change in abundance is a natural consequence of the present model since heavy ions are being preferentially accelerated into the primary current region where the soft X-ray-emitting plasma is predominantly located.

5.2. Parallel Ion Acceleration

The acceleration of the ions along the field lines is very much slower than the perpendicular acceleration, primarily because of the large scale lengths along the magnetic field compared with those across the fields lines. Figure 16 shows the parallel velocities of the He ions during the rise phase (left-hand side) and decay phase (right-hand side). On comparison with Figure 12, it is seen that there is very strong perpendicular heating at \( \omega_{pe} t \approx 360 \) from the corona down to the chromosphere, there is very little parallel heating. By \( \omega_{pe} t \approx 720 \) near the hard X-ray peak there are some energetic tails in the ion distribution. At this stage, the highest energy particles are located near the primary energy release region at \( x/\Delta \approx 200 \). Those ions with large negative velocities (i.e., upflowing ions) primarily reside in the primary current region while those with high positive velocities are primarily in the return current regions.

During the decay phase, the quasi-static electric field, while decreasing, is still present and continues to produce upward acceleration of ions. This continued acceleration is seen in the right-hand side of Figure 16 as the development of upflowing ion beams (i.e., where the bulk of the ions have negative velocities) at \( x/\Delta \approx 400 \). At lower altitudes, collisions tend to produce sufficient scattering in velocity space that there is only a heat flux upward rather than a well-defined beam.

5.3. Implications for Soft X-Ray Line Emissions

In order to compare the above characteristics of the ion acceleration with those inferred from soft X-ray line emissions, the velocity distributions of the He ions were convolved with the density of the soft X-ray plasma \( n_x \) as defined in § 4.2 to give a pseudo–line profile \( F_{He} \) similar to that of the soft X-ray line emissions, that is,

\[
F_{He} = \int dx \, dy \, n_x(x, y) f(v, x, y) \ .
\]  

(10)

Two different projections or viewing angles are presented. The first, which is shown in Figure 17, is for the case of an observer viewing the flare along the field lines, as for disk flares; here the Doppler shifts are solely due to the parallel velocity \( v_x \). The second case (Fig. 18) is for limb flares where the perpendicular velocity component is primarily responsible for the Doppler shifts. In each figure, the profile in Figure 17a is superposed to highlight the development of any broadening or blueshifts.

In Figure 17 it is seen that the theoretical line profile is initially dominated by a stationary component whose amplitude grows during the rise phase. As indicated in the previous section, the acceleration is relatively slow along the magnetic field so that only a relatively weak blueshifted component is seen during the rise phase. The maximum speed of this component is seen to increase on average during the rise phase. During the early part of the decay phase, the relative intensity of this blueshifted component is seen to increase as does its maximum speed, which reaches a peak value of about 900 km s\(^{-1}\) at about \( \omega_{pe} t \approx 1320 \). After this point, the maximum velocity of the blueshifted component is seen to decrease, and it is at this time that the break in the hard X-ray spectrum vanishes (Fig. 10). This behavior of the development of the blueshifted component relative to the stationary component is in agreement with the observations of disk flares presented by Wingee et al. (1991).

The main difference between the theoretical and the observed line profiles is that the stationary component decreases too fast during the decay phase. This difference is primarily due to the rapid loss of electrons from the flux tube, caused by the lack of a mirror force; thus, \( n_x \) decreases faster...
Fig. 17.—Calculated line profile for the He ions assuming the viewing angle is along the magnetic field. This case is important for comparison with soft X-ray line emissions from disk flares. A blueshifted component is seen to develop, particularly during the later part of the rise phase and early part of the decay phase where the fastest upward velocities are seen.
than it should. Since the ion dynamics are relatively slow in evolving, a more comprehensive treatment is likely to show only changes in absolute flux level, without any significant changes in profile, except possibly near the very end.

In comparison, the line profiles shown in Figure 18 which correspond to a cross-field line of sight show only a broadened stationary component. This component shows strong non-thermal broadening even at the earliest time (Fig. 18a) before the hard X-ray emissions begin and before strong blueshifts appear (e.g., dotted curve). This rapid broadening is due to the perpendicular acceleration of the ions over relatively short length scales (§ 5.1). The broadening reaches a maximum near

![Graphs showing velocity distribution over time.](image)

Fig. 18.—As in Fig. 17, except the viewing angle is orthogonal to the magnetic field. This case is important for comparison with soft X-ray line emissions from limb flares. Strong broadening is seen very early in the rise phase, and it reaches its maximum value near the hard X-ray peak.
the hard X-ray peak at $\omega_{pe} t \approx 840$ and then rapidly decreases except for a high-energy tail, consistent with the observations of limb flares presented by Winglee et al. (1991) and references therein.

6. CONCLUSIONS

In this paper, a particle simulation model for the acceleration and heating of electrons and ions during impulsive solar flares has been presented. The model extends previous work by including (1) a mechanism for the self-consistent acceleration of the hard X-ray-emitting electrons, (2) the development of the return current system in the corona and chromosphere, and (3) the outflow of chromospheric ions. The development of the current system is important since the associated electric fields can accelerate particles along the entire length of the flux tube, that is, the acceleration and heating is not restricted to the primary energy release region. Waves driven unstable by this particle acceleration can produce strong heating and high-energy tails in the particle distributions which have important consequences for the interpretation of X-ray and microwave emissions.

It should be noted the model uses compressed length and time scales in order to self-consistently treat the effects of plasma processes. The length and time scales are about $10^{-4}$ smaller than typically taken for flares. However, the ordering of plasma length and time scales is maintained with the duration of the simulations being comparable to a transit time of accelerated chromospheric ions out to the top of the loop (which in an impulsive flare is of the order of tens of seconds) and much longer than time scales for most of the relevant plasma processes. By maintaining this ordering, the potential experienced by the particles should be similar to that experienced in actual flares. Thus, the effects of wave-particle interactions and energy transport associated with the flow of particles along the field lines are incorporated in a consistent fashion. The simulations provide insight into a variety of processes (from flare initiation to the development of the flare current system and characteristics of the electron and ion acceleration) which to date have only been treated in a piecemeal fashion.

In the model the particle acceleration is initiated by a cross-field current in the corona. Such a current may arise, for example, from reconnection processes and/or current sheet acceleration. This current cannot be easily closed in the corona because of the limited cross-field mobility of particles in a collisionless plasma. As a result, the current system propagates down into the chromosphere where collisions can enhance the cross-field mobility and allow current closure. Thin current elements of the order of energetic ion gyroradii minimizes the distance that the charge carries have to travel across the field lines and so are preferred (see Winglee et al. 1988b).

The corresponding current densities and bulk acceleration of plasma along the flaring flux tube can be inferred from current closure and the geometry of the system (§ 3). The perpendicular current densities required to drive the system are relatively small compared to the parallel currents because of the aspect ratio of thin, long current elements. The electric field needed to support these currents produces bulk acceleration of electrons along the length of the flux tube, with high-energy electrons predominantly confined to the primary current region; the return current region, which is on adjacent field lines, remaining relatively cool. This electric field also accelerates ions across the field lines into the primary current region and then out into the corona.

The induced particle acceleration and heating and the associated X-ray and microwave emissions in this current system have the following characteristics:

1. Electrons in the primary current region are accelerated downward to hard X-ray energies by the quasi-static electric fields supporting the current system. Wave-particle interactions, particularly via Langmuir waves, scatter some of the electrons to energies at least 4 times higher. The combination of acceleration by the quasi-static electric fields and heating by wave-particle interactions produces an electron distribution with a broken-power law, similar to the observations of Lin & Schwartz (1987) and Winglee et al. (1991).

2. Within the primary current region, the ambient plasma also undergoes bulk heating to temperatures exceeding $10^7$ K (i.e., to soft X-ray-emitting temperatures). The extent of this hot plasma region propagates down into the chromosphere at about the ion sound speed, consistent with the soft X-ray fronts reported by Batchelor et al. (1985), Rust et al. (1985) and Nitta et al. (1989).

3. When this hot plasma reaches the chromosphere, the electron cross-field conductivity is enhanced by the large temperature gradient between the hot plasma and the surrounding, colder return current regions. As a result, the large electric fields initially required to close the current system are no longer needed, and this in turn limits the electron acceleration. As a result, the hard X-ray peak is reached at about this time.

4. Perpendicular heating of the electrons is produced via scattering from wave instabilities driven by $E \times B$ drifts. This energization occurs primarily for electrons in the beam region with moderate parallel velocities. As a result, the perpendicular electron energy and associated microwave emissions are expected to peak after the hard X-rays, which is consistent with the observed delay reported by Crannell et al. (1978), Kaufmann et al. (1983), and Cornell et al. (1984).

5. In contrast to the electrons, the ions first experience strong perpendicular acceleration into the primary current region to provide part of the perpendicular current needed to close the current system. Heavy ions are preferentially accelerated because of their large gyroradii, thereby increasing the heavy ion abundance in the primary current region. This increase in heavy ion abundance may account for those inferred by Sylwester et al. (1984) and Sylwester (1987) using soft X-ray line emissions.

6. This perpendicular energization may also account for the broadening of the stationary component seen in soft X-ray line emissions during limb flares. The broadening produced by this process is expected to begin before the main rise phase, reaching a maximum near the hard X-ray peak. During the decay phase, the broadening decreases because of the reduction in electric field strength as well as diffusion (particularly of the more energetic ions) from the primary current region into the cooler return current region. These characteristics are similar to those observed in the soft X-ray line emissions for limb flares (Winglee et al. 1991, and references therein).

7. The ions in the primary current region are accelerated out into the corona by the same electric field that accelerates the electrons down into the chromosphere. Because of their slow transit time, most of the ions experience continued acceleration through the rise and decay phases until the electric field vanishes, that is, when the break in the hard X-ray spectrum vanishes. This parallel acceleration can account for the correlation of the blueshifted component in the soft X-ray line emissions from disk flares with the hard X-ray spectrum (Winglee et al. 1991).
8. Those ions that remain in the return current region are accelerated down into the chromosphere; their energy on average is smaller than that of the ions in the primary current region. Since these induced ion flows reside in a current system, there is approximate current (or equivalently momentum) balance between the ion upflow in the primary current region and the ion downflow in the return current region (approximate because the electrons as well as ions carry current and momentum). Such momentum balance has been reported by Zarro et al. (1988), and it has been proposed as evidence for chromospheric evaporation. In the present model, this momentum balance arises because the particle acceleration and heating are occurring in a closed current system.

One important feature not included in the present model is a convergent magnetic field. During the rise phase, the electron pitch angles are sufficiently small that these effects are probably negligible. However, during the decay phase many of the electrons have experienced strong perpendicular heating and could be trapped in the corona. Indeed, in the present model the loss of soft X-ray emitting electrons is too fast. Mirror forces should also be important for the dynamics of the outflowing ions since they have large perpendicular energies, and some of this energy may be converted to parallel energy through conservation of adiabatic invariant as they flow out into decreasing magnetic fields. These effects are currently being investigated.

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APPENDIX

DETAILS OF SIMULATION MODEL

A.1. FIELD SOLUTION

The algorithms for the field solutions and particle pushers used in the simulations are as follows. At each time step, the charge density \( \rho \) is accumulated on the simulation grid using a linear interpolation of the particle charges onto the nearest grid points. This charge density is then used in Poisson's equation to evaluate the electric field, that is

\[
\nabla \cdot E = 4\pi \rho . 
\]

This equation in Fourier space reduces to

\[
E_x = -i4\pi \frac{k_x^2}{k^2} \rho_k , 
\]

\[
E_y = -i4\pi \frac{k_y^2}{k^2} \rho_k , 
\]

where \( k_x, k_y \) are the wavevectors in the x and y directions, respectively, and \( k (\neq 0) \) is the magnitude of the wavevector.

The resultant electric field, which contains contributions from both quasi-static electric fields and wave-particle interactions, is used in the equation of motion to determine the change in particle velocity, that is

\[
\frac{dv}{dt} = q \left( \frac{E}{c} \times B \right) , 
\]

where \( B \) is the magnetic field which is assumed to be uniform and in the x direction. This new velocity is then used to update the particle positions, and the process is repeated for subsequent time steps. Note that the equation of motion is dependent only on the charge-to-mass ratio \( q/m \) and not the individual values of \( q \) or \( m \). As a result, a plasma can be represented by superparticles with the same total charge and mass as the actual plasma but with their number being much smaller than in the actual plasma.

Periodic boundary conditions in the direction perpendicular to the magnetic field are assumed since it is envisioned that the actual flare consists of thousands or tens of thousands of the current elements considered here. The system is also assumed to be symmetric about both axes, which allows particles to be loaded into only one-quarter of the system and the charge and current densities in the other quadrants to be inferred from symmetry. The use of symmetry conditions provides significant savings in computer memory and time.

A.2. COLLISIONS

Across the transition region and chromosphere, the plasma density is sufficiently high and the plasma temperature sufficiently low that the plasma is collisionally dominated. Inclusion of these collisions is important since they provide an energy sink for the precipitating energetic electrons and a means for closing the currents induced in the corona by the flare acceleration process. The actual collisional processes involved vary with particle species and energy.

Important types of electron collisions incorporated into the simulations are (1) (low-energy) electron-electron elastic collisions, (2) (high-energy) beam-electron-chromospheric electron inelastic collisions, and (3) low-altitude, high-energy electron-neutral ionizing collisions. At the lowest altitudes, where the most energetic electrons are stopped (i.e., at about 1000 km), the scattering collision frequency for ambient energies of \( \leq 3 \) eV is about \( 3 \times 10^6 \) s\(^{-1} \), assuming a plasma density of about \( 10^{11} \) cm\(^{-3} \). At the higher
altitudes where the intermediate energy (10–50 keV) electrons stop (i.e., at several thousand kilometers), the plasma density and hence collision frequency are 1/10 as large. These collision frequencies are very much smaller than the local electron cyclotron frequency ($\Omega_e/2\pi \approx 1$ GHz) but increase from below the ion cyclotron frequency ($\Omega_i/2\pi \approx 0.5$ MHz) at several thousand kilometers to above the cyclotron frequency near an altitude of $\sim 1000$ km.

The dominant collisions for the ambient ions are electron-ion collisions and ion-neutral scattering collisions, the latter primarily important at the lowest altitudes (ion-ion collisions are relatively slow and not considered here). The cross section for these two collisional processes are approximately a factor of $(m_i/m_e)^{1/2}$ smaller than the electron-electron and electron-neutral collisions, respectively. Thus, the ion collision frequency ranges from a few percent of $\Omega_i$ to about $\Omega_i$. The perpendicular conductivity produced by these collisions is $(m_i/m_e)^{1/2}$ larger than that of the electrons, assuming $v_e \approx (m_i/m_e)^{1/2} v_i \ll \Omega_i$ (Krall & Trivelpiece 1973).

The above collisional processes are included in the simulation as follows. Electrons in the transition region and chromosphere are binned in space and energy. A total of 20 spatial bins (equivalent to 24$\Delta$ in $x$) and five energy bins (with a difference in speed between bins of $2v_T$) are used. The electron-electron collision frequency is assumed to have the form

$$v_e = v_{e0}(N_e + 1)^{1.4}(N_e)\frac{1}{1} ,$$  \hspace{1cm} (A4)

where $N_x$ is the bin number starting from the beginning of the transition region, $N_e$ is the energy bin number, and $v_{e0}$ is the initial electron collision frequency, and the ion collision frequency is assumed to be $(m_i/m_e)^{1/2}$ smaller. The collision frequency given by equation (A4) incorporates increases in the collision frequency with decreasing altitude, and decreases in the collision frequency with incident electron energy. These dependences on altitude and energy incorporate the basic physical processes, although it should be noted that the energy dependence is lower than actual Coulomb collisions, primarily to ensure that the bulk of the energetic electrons are stopped before they reach the simulation boundaries.

In the simulations, $v_{e0}$ is taken to be $0.2\Omega_i$, so that $v_e$ runs from about $0.5\Omega_i$ to $15\Omega_i$ for the electrons in the lowest energy bin. While the dynamic range (which is limited by numerical considerations) is smaller than in the actual chromosphere, test simulations with variations in collision frequencies and chromospheric densities by factors of 4 did not show significant differences.

Once the collision frequency is prescribed, the total number of electrons in each bin experiencing collisions is then determined at each time step. Those electrons in the lowest energy bin are assumed to experience only elastic scattering collisions. This assumes that energy losses from the ambient plasma can be neglected. The width of the first energy bin has been selected so that the ambient plasma can be heated to preflare coronal temperatures before energy losses become significant (cf. Winglee 1989). The electrons experiencing collisions in the other energy bins have their velocity reduced to two-thirds their initial value and scattered according to the differential cross sections given in Mott & Massey (1965). This degradation of incident velocity provides an energy sink for the precipitating energetic electrons (e.g., via radiative transfer and Coulomb collisions). No increase in the velocity of the target ambient electron is explicitly included, so that the present method tends to underestimate the production of mildly energetic electrons. However, intermediate energy electrons with their lower energy and high collision frequency are expected to lose their energy before they reach the simulation boundaries.

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While it is clear that the above model makes significant approximations to the variety of collisional processes involved in the dynamics of the chromosphere, we believe that it is an important first step to understanding the relative roles of the coronal and chromospheric plasma in determining the overall flare dynamics. The model presented here is one of the few models that treats particle acceleration in both the collisionless environment of the corona and the collisional environment of the chromosphere.

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