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ON THE IMPORTANCE OF REVERSE CURRENT OHMIC LOSSES IN ELECTRON-HEATED SOLAR FLARE ATMOSPHERES

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A. Gordon Emslie

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ABSTRACT

We consider the passage of a beam of non-thermal electrons through the flaring solar atmosphere, paying particular attention to the requirement that the beam be stable to the generation of plasma turbulence. We then compute the ratio of energy losses due to reverse current ohmic heating, and heating by Coulomb collisions, respectively, for the greatest flux which can pass stably through the atmosphere. We show that this ratio is determined by the low energy cutoff of the beam, by the electron temperature of the ambient atmosphere, and by the electron to ion temperature ratio θ . It is also independent of the atmospheric density. The results show that ohmic energy losses are undoubtedly important in the initial transient state, in agreement with other authors, but that their role is debatable in the flare atmosphere, depending on the value of θ appropriate. Expected values for θ during the impulsive phase of the flare indicate that reverse current ohmic energy losses, and their consequent effects on the electron beam dynamics and the hard X-ray bremsstrahlung emission, may not be as important as previously suggested; however, a fully time-dependent analysis of the beam-target interaction is necessary to fully resolve the issue.

Subject headings: hydromagnetics - particle acceleration - Sun: corona

Sun: flares - Sun: X-rays

I. INTRODUCTION

The necessity for a beam-neutralizing reverse current in a thick target electron-heated scenario of solar flares is by now well established (e.g. Hoyng, Brown, and van Beek 1976; Brown and Melrose 1977; Knight and Sturrock 1977; Hoyng, Knight, and Spicer 1978; Emslie 1980). In addition, detailed knowledge of the various processes affecting the passage of an electron beam through the solar atmosphere, and their relative importance, is essential in order to infer the characteristics of the accelerated electron population (and so place constraints on the flare primary energy release process) from radiation signatures such as hard X-ray emission. For this reason it is important to ascertain as accurately as possible the effect of the reverse current on the evolution of the distribution of accelerated non-thermal electrons.

A quantitative treatment of the ohmic heating of the flare atmosphere resulting from the passage of the reverse current through the resistive ambient plasma was performed by Emslie (1980). He showed that these ohmic losses can, under certain conditions, result in considerable modification to the dynamics of the electron beam (compared to those calculated under a purely collisional treatment - see Emslie 1978). In addition, he showed how the features (e.g., intensity, spectrum) of the hard X-ray bremsstrahlung produced by collisions of the beam electrons on ambient protons may also be significantly affected. In his analysis, however, he assumed, without rigorous justification, that the reverse current drift velocity was such that plasma instabilities did not develop. Hoyng, Knight, and Spicer (1978) did consider the question of the stability of the reverse current (against the generation of ion-acoustic turbulence) and concluded that if a flux F_{unstab} of non-thermal electrons is injected with an associated reverse current which would be unstable, then the instability acts to reduce this injected flux to a marginally stable value $F_{crit} < F_{unstab}$ (see also Brown and Melrose 1977; Manheimer 1977; Spicer 1977).

It is therefore of interest to consider the relative roles of ohmic (i.e., due to the passage of the [low velocity] reverse current through the finite resistivity background plasma) and collisional (i.e., due to the interactions of the [high velocity] beam electrons with the ambient particles) heating for a beam of non-thermal electrons whose flux, at injection, is such that the beam is marginally stable to the generation of plasma turbulence. We find (§II) that the ratio of the energy loss rates by both processes is dependent on the low energy cutoff in the injected electron spectrum, on the electron temperature of the target, and on θ , the ratio of electron (T_e) to ion (T_i) temperatures in the target. The first two of these parameters are quite well established (to within factors of two); however, the value of θ is not so well established, and the results are found to be quite sensitive to this parameter. We thus find (§ III) that ohmic energy losses will only be important if θ is comparatively low ($\lesssim$ 3), while for larger values of θ a purely collisional treatment is adequate. This threshold value of θ is then compared with those expected under a variety of conditions, allowing us to draw qualitative conclusions regarding the role of reverse current ohmic heating at various stages in the flare.

II. THE RATIO OF OHMIC TO COLLISIONAL ENERGY LOSSES FOR A MARGINALLY STABLE ELECTRON FLUX

We shall consider the injection of a beam of non-thermal electrons with energy spectrum

$$F_o(E_o) = (\delta-1)\,\frac{F_1}{E_1}\left(\frac{E_o}{E_1}\right)^{-\delta} \text{ electrons cm}^{-2}\text{s}^{-1}\text{keV}^{-1} , \qquad (1)$$

for $E_o \geq E_1$, the low energy cutoff (keV) for the beam. (Note that F_1 is thus the total injected flux.) Ohmic energy losses are greatest near the acceleration site (see Figure 3 of Emslie 1980) and so we will compare ohmic and collisional energy loss rates at that point. Using equation (34) of Emslie (1978) to calculate the collisional losses (assuming a fully ionized target) and setting the reverse current losses equal to $\eta j^2 = \eta e^2 F_1^2$ (η = plasma resistivity, j = reverse current density), we find that the ratio of ohmic to collisional losses is given by

$$\beta = \frac{\delta \eta F_1 E_1}{2\pi e^2(\delta-1)\Lambda n} , \qquad (2)$$

where e is the electronic charge (e.s.u.), n the target density (cm^{-3}), and Λ the Coulomb logarithm (Spitzer 1962).

Since the electron flux produced by any source will be limited to the marginally stable value, as discussed by Hoyng, Knight, and Spicer (1978; see discussion in §I), we may obtain an upper limit to β , β_o , by setting $F_1 = F_{crit}$, the flux at marginal stability. Further, since outside the electron source, i.e. in the region of beam propagation, the flux will always be lower than this (due to collisional and reverse current ohmic attenuation processes [see Emslie 1980]), we may set all plasma transport coefficients to their classical values (Spitzer 1962). In particular, therefore, the resistivity η will assume the form

$$\eta = \frac{\pi^{3/2} m_e^{1/2} e^2 \Lambda}{2\left(2kT_e\right)^{3/2}} \quad , \tag{3}$$

where m_e is the electron mass and k is Boltzmann's constant. F_{crit} is given, for marginal stability, by

$$F_{crit} = \alpha n v_e \ , \tag{4}$$

where $v_e = (kT_e/m_e)^{1/2}$ is the electron thermal velocity and α is a function of the electron to ion temperature ratio θ . Substitution of equations (3) and (4) into equation (2) yields an upper limit to the contribution of ohmic losses (relative to collisional ones) in the target:

$$\beta < \beta_0 = \left(\frac{\pi}{128}\right)^{\frac{1}{2}} \frac{\delta}{(\delta-1)} \left(\frac{E_1}{kT_e}\right) \alpha \quad . \tag{5}$$

Note that this expression is independent of the background density n .

The form of $\alpha(\theta)$ depends on the conditions in the region where the electrons stream. Duijveman, Hoyng, and Ionson (1981) have considered the effect of both ion-cyclotron and ion-acoustic wave generation on a driven electron flux. They find that for $\theta \gtrsim 8$ the marginally stable flux is determined by the ion-acoustic turbulence threshold; the relevant form is then as in Figure 8 of Fried and Gould (1961)[1]. For

[1] Some authors (e.g., Brown, Hayward, and Spicer 1981) have used an analytic approximation for $\alpha(\theta)$, which scales as $(1 + A\theta^{3/2} \exp[-\theta/2])$, where $A = (m_p/m_e)^{1/2} e^{-3/2}$, m_p being the proton mass. This function in fact fits the exact numerical results of Fried and Gould (1961) only for $\theta \gtrsim 6$. It also exhibits a maximum turning value of ≈ 0.28 at $\theta = 3$; there is clearly no justification for use of this maximum value, especially when one considers that it lies well outside the range of applicability of the approximation.

$\theta \lesssim 8$, however, the situation is somewhat more involved. For such low values of θ the onset of ion-cyclotron turbulence precedes the onset of ion-acoustic turbulence (i.e. the critical drift velocity for the former is smaller; see Figure 1A of Duijveman, Hoyng, and Ionson 1981); which stability threshold is relevant then depends on the turbulence level at which the ion-cyclotron waves saturate. If the ion-cyclotron wave turbulence does saturate (i.e. the saturation level is sufficiently low), then one may drive electrons to velocities larger than the ion-cyclotron stability threshold; α is then determined by the ion-acoustic stability threshold, as for $\theta \gtrsim 8$. If, on the other hand, the waves do not saturate (due to a high saturation level), then the ion-cyclotron stability threshold determines α . Since the saturation level of ion-cyclotron waves in solar conditions is not well known (Duijveman, Hoyng, and Ionson 1981), it is in practice uncertain which of these two possibilities for $\theta \lesssim 8$ prevails.

Figure 1 In Figure 1 we show the dependence of β_0 on θ , determined from equation (5) and the above considerations of the value of α appropriate, for a typical $\delta = 4$ and for various values of E_1/T_7 , where E_1 is measured in keV and T_7 is the coronal electron temperature in units of 10^7 K . The solid lines correspond to the ion-acoustic threshold and the dashed lines to the (possibly irrelevant, depending on the turbulence saturation level) ion-cyclotron threshold. Note that when $\theta \lesssim 8$ and the ion-cyclotron wave saturation level is low, the resistivity η_{IC}, associated with the saturated ion-cyclotron turbulence, is less than the Spitzer (1962) classical value (see Figure 5 of Duijveman, Hoyng, and Ionson 1981). Thus classical resistivity remains a valid approximation, so that equation (5), and also the solid curve in Figure 1, still give a good estimate of β_o under these conditions.

III. DISCUSSION

In order to interpret the results of Figure 1, we must first assign an appropriate value to E_1/T_7 . T_7 increases from ≈ 0.3 in the preflare state (Noyes 1971) to a maximum of around (2 ± 1) (e.g. Datlowe, Hudson, and Peterson 1974). Power-law X-ray spectra have been observed down to $\varepsilon \approx 5$ keV (Kane et al. 1979); when one takes into account the thermal contribution to the photon flux at these energies, one finds that this corresponds to $E_1 \approx 7 - 8$ keV (Brown, Hayward, and Spicer 1981). We shall therefore adopt $E_1/T_7 \approx 25$ for the preflare state, and $E_1/T_7 \approx 10$ as an upper limit for the flare itself.

Considering first the preflare atmosphere, reference to Figure 1 shows that β_0 exceeds unity for all values of $\theta \lesssim 5$, whatever the marginally stable state in the target may be. Thus, since $\theta \approx 1$ in the preflare state, reverse current ohmic energy losses can easily dominate over collisional losses, confirming the conclusions of Knight and Sturrock (1977). However, in the flare atmosphere, β_0 exceeds unity only for $\theta \lesssim 3$ in the case of an ion-acoustic turbulent threshold, and only for $\theta \approx 1$ for an ion-cyclotron turbulent threshold.[2] Because

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[2] The results of Emslie (1980), showing that reverse current effects are strong for large injected fluxes, tacitly assume that such large fluxes are stable to the generation of plasma turbulence. This implies that a reverse current with a large drift velocity be permitted to pass stably, which in turn implies that θ be close to unity (Fried and Gould 1961; Duijveman, Hoyng, and Ionson 1981).

Coulomb collisions and ohmic heating both preferentially heat electrons as opposed to ions, it would appear that, since the initial ion temperature is only of order 10^6 K, the appropriate value of θ is much larger than either of these two values, implying $\beta_0 \ll 1$.

However, the heating of the plasma by both collisional and ohmic processes is counteracted by conductive and radiative cooling of the heated region. If a steady state situation obtains, then the external source terms in the equations controlling the behavior of T_e and T_i will vanish and T_e and T_i will be governed simply by the equilibration equations

$$\frac{dT_e}{dt} = \frac{T_i - T_e}{\tau} \quad ; \quad \frac{dT_i}{dt} = \frac{T_e - T_i}{\tau} \quad , \tag{6}$$

where

$$\tau = 12.6 \, T_e^{3/2} \, n^{-1} \text{ seconds} \tag{7}$$

is the (classical; see §II) electron-ion temperature equilibration time (Spitzer 1962). The solution of equations (6) and (7) has been obtained by Sivukhin (1966, his equation [9.17] and Figure 6); for an initial state in which θ is large (due to preferential heating of electrons; see remarks above) one finds that θ falls to a value $\lesssim 3$ in a time $\tau^* \approx 0.3\, \tau_o$ and to a value ≈ 1 in $\tau^* \approx \tau_0$, where τ_o is the initial temperature equilibration timescale. The first of these τ^* values corresponds to $\beta_0 = 1$ when the ion-acoustic turbulent threshold applies and the latter to $\beta_o = 1$ when the ion-cyclotron turbulent threshold applies (Figure 1).

We thus see that if the beam injection time is long compared to τ^* and to the time taken to achieve a balance between source and sink terms in the energy equation, then a state of equal electron and ion temperatures will be reached, and reverse current ohmic energy losses can in fact become important again, at least near the electron injection site (Figure 1; Figure 3 of Emslie 1980). If, on the other hand, either τ^* or the time to achieve local energy balance is much larger than the duration of the electron input, then θ will remain large and so reverse current ohmic energy losses will be unimportant (relative to collisional ones) everywhere in the target.

These considerations reveal that the comparative roles of reverse current ohmic, and direct collisional, energy losses are determined by the target density, not (as might be expected) directly through its appearance in the collisional energy deposition rate (erg cm^{-3} s^{-1}), but through its effect on the local energy balance in the flaring corona and on the electron-ion temperature equilibration time τ (eq. [7]). Typical preflare coronal densities are $\gtrsim 10^{10}$ cm^{-3} so that (assuming hydrodynamic effects are small) $\tau_0 \lesssim 40$ s. We thus see that τ^* could be as low as 10s (when the ion-acoustic turbulent threshold applies); this is comparable to the duration of an "Elementary Flare Burst" (de Jager and de Jonge 1978), allowing significant temperature equilibration to occur should balance between heating and cooling terms be achieved within such a timescale. If, however, a state of local energy balance takes longer than this to be attained, then θ will remain $\gg 1$ and so β_0 will remain $\ll 1$, as discussed above.

In summary, therefore, the relative roles of reverse current ohmic, and collisional heating are determined by the magnitude of the injected non-thermal electron flux, in addition to the parameters of the target atmosphere. An upper limit to this injected flux is that which is marginally stable to the generation of plasma turbulence, either ion-acoustic or ion-cyclotron as appropriate. In a preflare atmosphere, an injected flux of this magnitude produces strong ohmic heating, in agreement with the conclusions of Knight and Sturrock (1977). In a flare atmosphere, however, the importance of ohmic energy losses is determined principally by the critical reverse current drift velocity and hence by the electron to ion temperature ratio θ . This in turn depends on the relative sizes of three timescales: the electron beam lifetime, and the timescales for achieving local energy balance, and equilibration of electron and ion temperatures, respectively. If a state of local energy balance is reached while the electrons are still being injected, then θ will approach unity within the order of an electron-ion temperature equilibration time, after which reverse current ohmic energy losses will be an important consideration, at least near the electron injection point, for suitably large (but stable) injected electron fluxes (Emslie 1980). On the other hand, if a state of local energy balance cannot be reached this quickly, then θ will remain large, so rendering reverse current ohmic energy losses negligible compared to collisional losses throughout the entire atmosphere for any stable electron flux. A detailed study of the time-dependent response of the solar corona to a non-thermal electron energy input is necessary in order to resolve these questions. In addition, it is important to assess whether different "Elementary Flare Bursts" (corresponding to different "spikes" in the hard X-ray flux-versus-time profile -- see de Jager and de Jonge 1978) result from repeated re-energization of the same region, or to successive energizations of different regions (see Karpen, Crannell, and Frost 1979; Emslie 1981). In the former case, the time available for electron-ion temperature equilibration will be relatively large, while in the latter case each burst must be considered individually, with preflare initial conditions, thus requiring that the electron-ion temperature equilibration time be less than the duration of the Elementary Flare Burst if reverse current ohmic losses are to become an important energetic consideration at some stage in the burst.

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FIGURE CAPTION

Fig. 1 - Ratio of reverse current ohmic heating to collisional energy deposition, β_0 , as a function of the electron-ion temperature ratio θ . The values of β_0 have been calculated for a non-thermal electron flux whose associated reverse current, at the point of injection, is marginally stable to the generation of plasma turbulence. They are also evaluated at this injection point, and so represent an upper limit to the importance of reverse current ohmic heating in the atmospheric energy balance (Emslie 1980). The dashed curves correspond to the marginally stable threshold for ion-cyclotron turbulence, and the solid curves to the marginally stable threshold for ion-acoustic turbulence, applicable either when $\theta \gtrsim 8$ or when $\theta \lesssim 8$ and the ion-cyclotron wave saturation level is low (Duijveman, Hoyng, and Ionson 1981). In both cases β_0 depends linearly on the ratio of the low energy cutoff in the electron spectrum E_1 (keV) and the electron temperature of the target atmosphere (T_7 in units of 10^7K), as shown.

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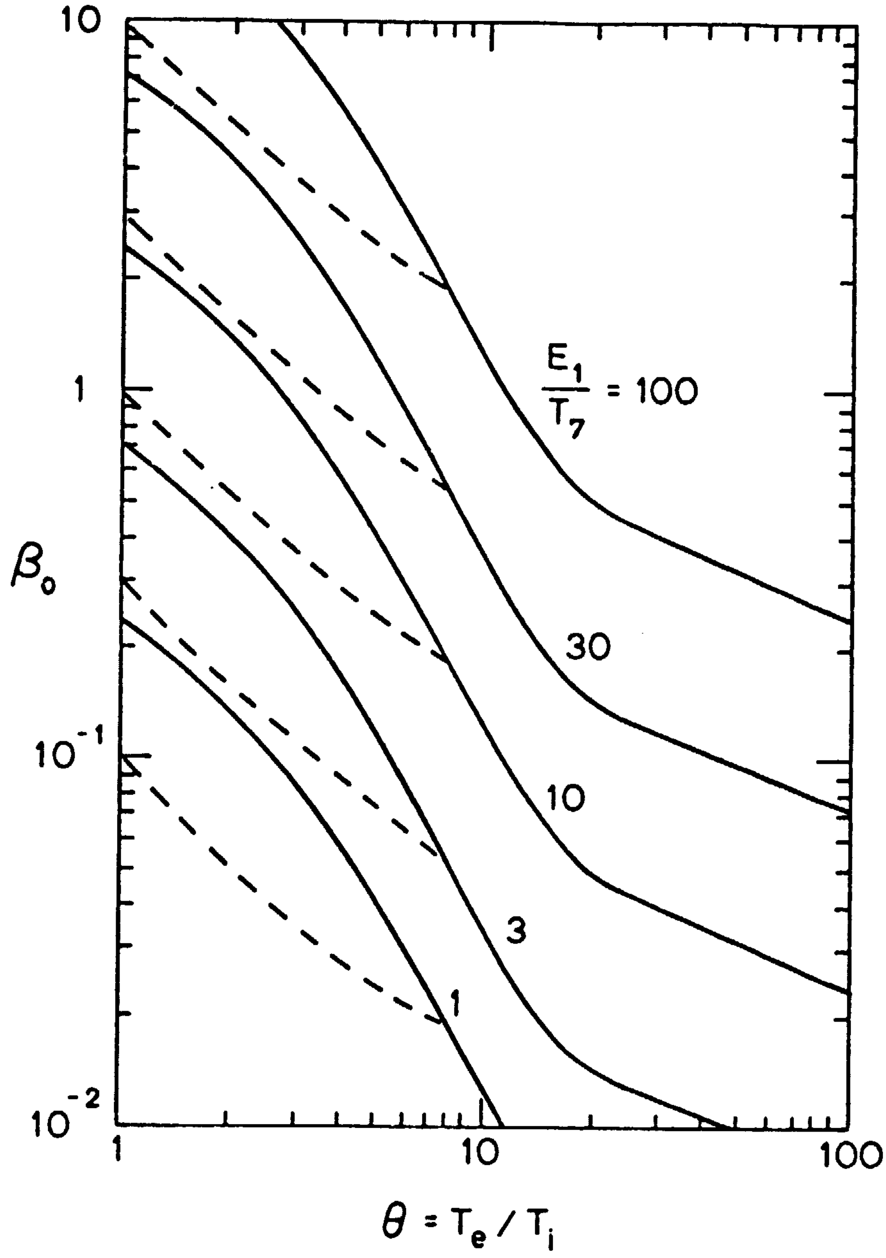

Figure 1

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20. ABSTRACT (Continue on reverse side if necessary and identify by block number)

The maximum flux of non-thermal electrons which can pass stably through a plasma is determined by the threshold for onset of plasma turbulence. Since the importance of reverse current ohmic losses (relative to beam-target collisions) in controlling the dynamics of such non-thermal electron beams in solar flares increases with the injected flux, we therefore find that the extent to which reverse current ohmic dissipation may be important is determined principally by the critical drift velocity of the reverse current, and

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so by the ratio of electron to ion temperatures in the target. The expected behavior of this ratio as a function of time is briefly discussed, with a view to ascertaining the importance of the reverse current in controlling the dynamics of the electron beam.