STOCHASTIC ACCELERATION IN THE DIFFUSION REGION AND THE STRUCTURE OF SLOW SHOCKS IN SOLAR FLARES

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Abstract

Stochastic particle accelerations in the diffusion region of magnetic neutral sheet expected in solar flares are studied. The bulk energy and short time scale of flares are found to be accounted for by electron accelerations due to the turbulent Lower Hybrid Drift waves. Independent from above, the Petschek-type slow shock, if it exists, is found to be collisionless, and the relevant energy conversion is discussed. Finally a comparison between the stagnant diffusion model and slow shock model is made.

1. Introduction

Yohkoh has revealed that the top of the geometrically larger soft X-ray flares shows a cusp shape, suggesting the magnetic reconnection is occurring above it as conjectured by many authors, e.g. Hirayama (1974). In a recent short note (Hirayama, 1994, paper I), we investigated the main energy conversion in a purely antiparallel magnetic field configuration above the flaring loop. Unlike the popular view, this diffusion model (height extension of 10^4km and width of a few meters) was found to be able to convert the magnetic energy very rapidly because of the anomalous resistivity from possible generation of lower hybrid drift (LHD) instability. It is found to be imperative, however, that most of the energy be converted to high speed particles, which can only escape from the same side of 10^4km long where the magnetic energy was brought in. This is an essential point for the diffusion model to survive. Therefore we first inspect stochastic particle accelerations due to LHD- waves in the stagnant diffusion region of Paper I. Secondly, independent from the diffusion model, the structure of the possible Petschek-type slow shock is studied, and finally we compare the two models. We take antiparallel magnetic fields **B** in the vertical, z-direction, and x-axis parallel to the solar surface, and $\partial/\partial y = 0 = \partial/\partial t$ is assumed. We use cgs-Gaussian unit ($e = 4.8 \cdot 10^{-10}$, emu unit in Paper I), temperature in energy unit (only) in equations (no Boltzmann's constant is used), and also thermal velocity in $(2T/m)^{1/2}$ ($(kT/m)^{1/2}$ in Paper I), and suffix i and e for ions (protons) and electrons. The lower hybrid frequency is nearly equal to $\Omega_{LH}(\mathrm{s}^{-1}) \equiv eB/c\sqrt{m_i m_e} = 4.1 \times 10^7 (B/100\mathrm{G})$.

2. Particle Acceleration due to LHD Waves in the Diffusion Region

Once the cross-**B** drift velocity of the electric current, v_d ($=J_y/en$), approaches ion thermal velocity due to thinning of the neutral sheet, the LHD plane waves propagating in the y-direction are generated with wide frequency ranges centered at $\Omega_{\rm LH}$ (e.g. Brackbill *et al.*, 1984; Davidson *et al.*, 1977). We consider the LHD-waves primarily because they are excited for smaller v_d than other plasma waves in the preflare corona, where $T_e \approx T_i$ and $T_\perp \approx T_\parallel$ are expected (e.g. Gary, 1993). The particles are, however, essentially bound in a gyroradius. Drift velocities do exist as a cumulative effect of spatial differences of a number of gyrating particles due to density gradient (Spitzer, 1962, p.32). The amplitudes of particle displacements oscillating with fluctuating electric field $\delta E_y (\gg \delta E_x; \delta B_z \approx 0)$ are also small, at most an electron gyroradius, a_e. The recognition of bound particles is the basic ingredient of this paper. Then particles are stochastically accelerated by δE_y, almost in situ except for the anomalous diffusion discussed below.

We define a time span during which $\delta E_y(t)$ holds a constant value as τ_{cor} (autocorrelation time $\equiv 1/\Delta\Omega_{\rm LH} \equiv Q/\Omega_{\rm LH}$). Then particles (mass m, charge q), oscillating between $\Omega_{\rm LH} - \Delta\Omega_{\rm LH}$ and $\Omega_{\rm LH} + \Delta\Omega_{\rm LH}$, gain or lose momentum by $\Delta p_y = m\Delta v_y = q\delta E_y \tau_{cor}$ through a single encounter with a (constant) δE_y during τ_{cor}. For a longer time Δt, there should be $N = \Delta t/\tau_{cor}$ encounters, of which a number of $\sqrt{N}$ will act as an acceleration due to random walk: $\Delta p = m\Delta v = q\delta E_y \tau_{cor}\sqrt{N}$. Here Δv and δE_y are taken to be *rms*-values. The average rate of the energy gain per particle is then $dE/dt = m < (\Delta v)^2 > /2\Delta t = (e\delta E_y)^2 \tau_{cor}/2m$. This coincides with formal Fourier treatments.

The number density of particles n_E between the energy E and $E + \Delta E$ will be governed by $\partial(n_E dE/dt)/\partial E = -n_E/\tau_{esc}$ with $\tau_{esc} \equiv v_{esc}^{-1}\partial x/\partial \ln(n_E)$, a characteristic escape time. This equation for n_E comes from the collisionless Boltzmann (or Liouville) equation for the momentum distribution n_p, $\partial(n_p dp/dt)/\partial p + \partial(n_p dx/dt)/\partial x = 0$, where dp/dt and dx/dt are ensemble averages (Braginskii, 1965, eqn. (1.1) and below), and n_p is converted to n_E using $n_p dp = n_E dE$ and $dE/dp = \sqrt{2E/m}$. The solution, using the above $dE/dt = \frac{1}{2} m_i v_d^2 Q \Omega_{\rm LH}(m_e/m)$ for $\delta E_y = B_z v_d/c$ is simply $n_E = (n_f/E_0)\exp(-E/E_0)$. We used the saturated field of $\delta E_y = k_y \delta\phi = B_z v_d/c$ from the saturated potential $\delta\phi = (T_i/e) \cdot (5/3)(m_e/m_i)^{1/2} v_d/v_{Ti}$ in Brackbill *et al.* (1984) from many simulations, noting $a_e k_y (T_i/T_e)^{1/2} \approx 1$, where k_y is the wave number ($\delta E_y^2/8\pi n_i T_i = (v_A/c)^2 (v_d/v_{Ti})^2 \ll 1$). Here $n_f({\rm cm}^{-3}) = \int_0^\infty n_E dE$ (suffix f refers to $fast$ particles), and $E_0 \equiv \frac{1}{2} m_i v_d^2 Q^2 \cdot (\tau_{esc}/\tau_{cor})$ for electrons (ion mass enters here). The E_0 for protons is by a factor m_i/m_e smaller, and thus they are practically not accelerated. Power law energy spectra may be obtained by summing different exponential distributions thus obtained, or by considering that particles of smaller $E_\parallel$ ($E_\parallel \ll E_\perp$ in our model) may be thermalized in the corona, or else by considering e.g. $\tau_{\rm esc}$ to be dependent on the particle momentum: $\tau_{\rm esc} \sim p^2$ does the job.

We derive the anomalous spatial diffusion related to $\tau_{\rm esc}$. For each encounter with the waves during $\tau_{\rm cor}$, particles will drift towards x-direction by $\Delta x = (\delta v_{xe})\tau_{\rm cor} = (c\delta E_y/B)\tau_{\rm cor} \approx v_d \tau_{\rm cor}$ due to the electric field drift of $\delta \mathbf{v}_E = c\delta\mathbf{E} \times \mathbf{B}/B^2$. For $N = \Delta t/\tau_{\rm cor}$ encounters during Δt, the expected drift becomes $\delta_x = \Delta x\sqrt{N} = v_d(\tau_{\rm cor}\Delta t)^{1/2}$. Taking Δt to be the escape time with the escape distance of $\delta_{esc} = \delta_x$, we obtain $\tau_{\rm esc} = (\delta_{\rm esc}/v_d)^2/\tau_{\rm cor}$, and finally $v_{\rm esc} = \delta_{\rm esc}/\tau_{\rm esc} = v_d(\tau_{\rm esc}/\tau_{\rm cor})^{1/2}$. Again these are consistent with the Fourier treatmetns of the *spatial* diffusion coefficient : $D_s \equiv \delta^2/\Delta t = v_d^2 \tau_{\rm cor} = c^2(\delta E_y)^2 \tau_{\rm cor}/B^2 \approx {\rm O}(a_e v_{Ti})$. Unlike the acceleration, both electrons and ions diffuse out with the same velocity.

Now we consider three conservation laws in the outer edge of diffusion region (Paper I) to be combined with the above E_0, $\tau_{\rm esc}$ and $v_{\rm esc}$: a) the escape flow of fast particles should be supplied by the diffusive influx from the e̲x̲thermal region (suffix e̲x̲), $n_f v_{\rm esc} = n_{ex} v_{ex}$, ensuring the zero fluid velocity, b) the energy outflow of escaping fast particles be supplied by the incoming Poynting energy flux, $n_f v_{esc} E_0 = B_{ex}^2 v_{ex}/4\pi (E_0 =$ mean energy per electron), and c) the pressure balance is $2n_cT_c = B_{ex}^2/8\pi$ (suffix c, *center* of the diffusion layer).

As an example we consider the case when $n_f = n_c$ is realized, in which case a maximum external inflow v_{ex} is reached (normally $n_f < n_c$). We adopt $Q = 0.1$ (a broad spectrum of δE_y in accord with linear analyses), $B_{ex} = 10^{1.5}$G, and $n_{ex} = 10^9$ ($v_{A,ex} = 2200$ km s^{-1}), together with a simplifying assumption of $E_0 = T_c$ consistent with $n_f = n_c$. We then obtain the following: $v_{ex}(\equiv v_{in}$ of Paper I) $= 77$km s^{-1}, or an Alfvén Mach number of the whole flare process $M \equiv v_{ex}/v_{A,ex} = 0.035$ (consistent with observations), a typical electron energy of $E_0 = m_i v_{A,ex}^2 = 50$keV (a satisfactory value), $n_f = n_{ex}/4.0$, $v_{\rm esc} = \sqrt{2}Qv_{A,ex} = 310$ km s^{-1}, $\delta_{\rm esc} = 24$cm, $v_d = 3100$km s^{-1}, $N = 100, \tau_{\rm esc} = \tau_{\rm acc} \equiv E_0 dt/dE = 0.77\mu$s (an extremely rapid process), and $\delta E_y = 98$ V cm^{-1}. In this case the energy conversion is occurring only at the outermost edge of the whole diffusion region of 510cm. Hence we consider that the bulk central region is rather quiescent without energy inflow, nor strong turbulence (possibly $B_z \approx 0$), perhaps except in the initial stage. Ions might be accelerated by ion-cyclotron turbulence which would be excited afterwards due to $T_e \gg T_i$.

The above scenario is still at a very primitive stage, and we need to study further the basic acceleration process more closely, and variations of physical parameters along x-axis.

3. Structure of the Petschek-type Slow Shocks

We inspect now the structure of the possible slow shocks. In the usual Petschek model the inflow velocity v_{x0} in unit of the preshock Alfvén wave velocity is given $M \equiv v_{xc}/v_{A0} = 0.01 \sim 0.1$. Plasmas are ejected at an Alfvén velocity of the inflow region in the $\pm z$ direction, and also the gas pressure rises to $P_1 \approx B_{z0}^2/8\pi$ from $P_0 \approx 2$ dyn cm^{-2} of preshock regions. Since the density enhancement over the initial state is ≈ 2.5 with $\gamma = 5/3$ (Bazer and Erickson, 1959), the temperature is expected to become very high: $T_1 = 2 \times 10^8 (B_{z0}/10^2G)^2/(n_0/10^{9.5})$K. However in order to obtain high temperature, there must be dissipation mechanism, Joule or viscous heating, in a 'collisional shock' through binary collisions. And for the binary collision to occur, at least the entire region has to be (much) larger than the mean free paths $l_{m0} = 145T_{6.4}^2/n_{9.5}$km in the preshock, and $l_{m1} = 3.1 \times 10^5$km after the shock. On the other hand the whole width of the post-shock region '1' is estimated to be 10^4km $\times\ M \approx 100 - 1000$km. This is much smaller than the averaged value of l_{m0} and l_{m1}: e.g. $l_m = (l_{m0}l_{m1})^{1/2} = 6700$km. Hence we must conclude that if the slow shock exists, it must be collisionless.

In the following we inspect whether a collisionless slow mode shock can be generated and converts the energy rapidly. We consider a situation, where the curvature or bent of the magnetic field due to vertically ejected flows becomes so large that the particle drift velocity v_d associated with the high current density $J_y = cB_z/4\pi\delta = en_e v_d$ becomes comparable to the ion thermal velocity. Then the LHD-instability will set in even for $T_e \approx T_i$. The electrical resistivity η^* becomes anomalous in such a way that the effective collision frequency ν^*(s^{-1}) becomes on the order of the lower hybrid frequency of $\Omega_{LH} : \nu^* \approx 0.01\Omega_{LH}(v_d/v_{Ti})^2/\beta_i$ where $\beta_i \equiv n_iT_i/(B^2/8\pi)$ (Brackbill *et al.*, 1984). The effective particle mean-free-path is now $l_m \approx v_{Te}/\nu^* \approx 3 \times 10^9$cm s$^{-1}/8 \cdot 10^5s^{-1} = 0.04$km ($v_d \approx v_{Ti}$ is assumed), very much smaller than the width of the whole region, and the slow shock is allowed to form. The thickness of this high current region is $\delta \leq 80$m $\times\ (v_{Ti}/v_d)B_2/T_{6.5}^{1/2}n_{9.5}$. Then we can calculate the Joule dissipation $\eta^* J_y^2 = m_e n_e v_d^2 \nu^*$ (η^*(s) $= \nu^* m_e/n_e e^2$), and the expected energy conversion to thermal energy within a thickness δ (shock thickness) is $\eta J_y^2 \cdot \delta$(erg cm^{-2}s^{-1}). The Joule heating should be compared with the horizontally incoming magnetic energy flux of $v_{x0}B_{z0}^2/4\pi$ (=Poynting energy flux). The ratio of the two is found to be $0.3B_2^2(v_d/v_{Ti})^3/(n_{9.5}T_{6.4}^{1/2} \times (v_{xo}/60\text{km s}^{-1}))$. The Joule heating is effective only in the lower temperature portion inside the shock transition region. This means that a substantial fraction of the incoming energy flux can be converted into thermal energy (within an 80m thickness), consistent with the jump condition of the MHD slow shock. The inflow velocity $v_{xo} = 60$km s^{-1}($\approx 0.05v_{A0}$) is chosen in accordance with

Paper I. However, a crucial problem is whether the scale length of the field change becomes in fact as small as 80m. The argument above is analogous to the Joule heating of the diffusion region (Paper I). The latter may now be regarded just as a demonstration because most of the energy release may be performed in the form of statistical accelerations, if the treatment in the preceding section is correct. In the same way we can discuss the statistical acceleration in the shock region. The reason is that the transit time of plasma in the shock front 80m/60km $s^{-1} \approx$ 1ms is much slow as compared with the particle acceleration (=escape) time of $\approx 1\mu$s. Therefore the same logic with almost the same numerical values may be applied to the shock.

4. Comparison between the Slow Shock Model and Diffusion Model

The above collisionless slow shock model and the diffusion model of Paper I, both seem to be tenable at least in the present preliminary study. The diffusion model requires a very small width of a few meters (in x-direction) to be able to generate LHD-instability and a large height extension of order of 10^4km (in z) to be able to convert enough energy. Whether it is likely to form or not is intimately related to whether the Petschek-type reconnection occurs or not. The tearing mode instability might occur in the long diffusion region which is stably sandwiched by the two walls of antiparallel fields like coronal helmet streamers. Since more slender magnetic islands thus created are more unstable, the situation will not differ very much without tearing. We note also that the inclusion of shear ($B_y \neq 0$) does not make much difference in the diffusion model (Krall, 1977), but might jeopardize the creation of slow shocks, and strong curvature in the shock may act to stabilize LHD instability (see paper list in Brackbill *et al.*).

The width of the collisionless shock should also be very narrow as estimated above ($\approx$ 100m) with a large height extension of 10^4km. A question is how the very first reconnection starts. A necessary condition of the generation of the Petschek-type reconnection is to make reconnection at a "single" point of order of less than 10 meters within a 10^4km-scale of the *smoothly curved* magnetic field. If reconnections start at a few, or thousand points in a 10^4km length, it becomes similar to the diffusion model. Another question in the Petschek-type model is that when the flare is associated with a prominence eruption, the flaring starts after the prominence reaches a height of $3 - 10 \times 10^4$km (e.g. Miyazaki *et al.*, 1994). This suggests that the vertical neutral sheet may already have been formed between the prominence and the chromosphere, in favor of the diffusion model. This will become explosive when the thinning is enough to generate a very high electric current, namely large drift velocity. Finally, we remark that in *Yohkoh* movies we have not detected *as yet* any evidence of down flows just above the (reconnected) flaring loops. I appreciate discussions with Drs. S. Hinata and K. Shibata.

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