STOCHASTIC PARTICLE ACCELERATION IN FLARING STARS

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ABSTRACT

We consider the acceleration of electrons by the Fermi-Parker mechanism in a quasi-stationary turbulent plasma of dimension ℓ, mean magnetic field strength B, and mean number density n. The electrons suffer radiative and ionization losses ($\dot{p} = -\alpha p^2 - \gamma$) and have a scattering mean free path that increases linearly with their momentum. We give exact analytic solutions for the steady-state electron energy spectra. The spectra are characterized by an exponential cutoff above a given momentum determined by the synchrotron or the confinement time, depending on the physical characteristics (ℓ, B, n) of the accelerating region.

1. Introduction. Direct acceleration of electrons from a thermal pool to relativistic energies is problematic for two reasons both, in diffusive shock wave acceleration, and in resonant Alfvén wave scattering. First, because of their small Larmor radii electrons do not see a shock wave as a discontinuity, unless their momentum becomes very large. Secondly, electrons only fulfil the resonance condition with Alfvén waves if their momentum is larger than $p_t = m_p v_A = 20$ keV/c $(B/3 \cdot 10^{-6}\ \mathrm{G})(n/1\ \mathrm{cm}^{-3})^{-1/2}$ (Achterberg and Norman 1980). One solution to this problem is to find injectors of energetic electrons in space which preaccelerate these particles to moderately large momenta $p > p_t$. Viable injector candidates are (i) secondary electron production in inelastic nuclear collisions of cosmic ray nucleons with interstellar gas atoms and molecules (e.g. Schlickeiser 1982), and (ii) electron acceleration in the flares on M and K type stars (Lovell 1974). Here we discuss the acceleration of energetic electrons by the second-order Fermi, or stochastic acceleration, mechanism in flare stars in more detail.

2. Acceleration in Flare Stars. We consider the acceleration of relativistic electrons by the Fermi mechanism in a quasi-stationary turbulent plasma of dimension ℓ, mean magnetic field strength B, and mean number density n. Additionally, we allow the electrons to suffer simultaneously radiation and ionization losses,

$$\dot{p} = -\alpha p^2 - \gamma \qquad (1a),$$

with

$$\alpha = 1.2 \cdot 10^{-11}\ (B/B_2)^2\ (\mathrm{eV/c})^{-1}\ \mathrm{s}^{-1} \qquad (1b),$$

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*The National Center for Atmospheric Research is sponsored by the National Science Foundation

$$\gamma = 5.2 \cdot 10^2 \; (n/n_9) \; (eV/c) \; s^{-1} \qquad (1c),$$

where (B/B_2) is the magnetic field strength in units of 100 Gauss, (n/n_9) the plasma density in units of 10^9 cm^{-3}. The electrons may also escape from the acceleration region with a mean free path that increases linearly with momentum, so that the escape time is

$$T = \ell^2 / K_{\|}(p) = 1 / (Ap) \qquad (2a)$$

with

$$A = 5.8 \cdot 10^{-9} \left(\frac{\lambda}{\lambda_7}\right) \left(\frac{\ell_{10}}{\ell}\right)^2 \; (eV/c)^{-1} \; s^{-1} \qquad (2b),$$

where ℓ/ℓ_{10} is the dimension of the acceleration region in units of 10^{10} cm, and λ/λ_7 the scattering mean free path ($\lambda = 3 K_{\|}/v$) in units of 10^7 cm at $p = m_e c$. The second-order Fermi diffusion coefficient then is (Skilling 1975)

$$D_{pp} = \frac{V_A^{\,2} \, p^2}{9 \, K_{\|}(p)} = D \, p \qquad (3a)$$

with

$$D = (2.7 \cdot 10^5) \; (B/B_2)^2 \; (\lambda_7/\lambda) \; (n_9/n) \; (eV/c) \; s^{-1} \qquad (3b).$$

The equilibrium phase space density of electrons resulting from the combined effect of acceleration, radiation loss, ionization loss and escape is given by the solution of (Schlickeiser 1985)

$$\frac{\partial f}{\partial t} = \frac{1}{p^2} \frac{\partial}{\partial p} \left[p^2 \left(D_{pp} \frac{\partial f}{\partial p} - \dot{p} f\right)\right] - \frac{1}{T} f + S(p,t) \qquad (4),$$

where $S(p,t)$ denotes the source term.

We define the dimensionless momentum $x = p/p_c$, where

$$p_c = (D/\alpha)^{1/2} = 1.5 \cdot 10^8 \left(\frac{\lambda_7}{\lambda}\right)^{1/2} \left(\frac{n_9}{n}\right)^{1/2} \; (eV/c) \qquad (5)$$

and the parameters

$$a \equiv \gamma/D = 1.9 \cdot 10^{-3} \; (n/n_9)^2 (\lambda/\lambda_7)(B/B_2)^{-2} = \frac{\tau_F(p)}{\tau_I(p)} \qquad (6)$$

$$b \equiv A/\alpha = 4.8 \cdot 10^2 \; (\ell/\ell_{10})^{-2} \, (\lambda/\lambda_7) \, (B/B_2)^{-2} = \frac{\tau_R(p)}{T(p)} \qquad (7).$$

p_c is that momentum where the Fermi acceleration time ($\tau_R = p/D$) exactly equals the radiation loss time ($\tau_R = (\alpha p)^{-1}$). a is the momentum-independent ratio of Fermi acceleration time ($\tau_F = p/D$) to ionization loss time ($\tau_I = p/\gamma$). b is the momentum-independent ratio of radiation loss time ($\tau_R = (\alpha p)^{-1}$) to escape time ($T = (Ap)^{-1}$).

3. Results and Discussion. We consider the delta-function source term ($x_o = p_o/p_c$)

$$S(x) \equiv \frac{N_o}{\tau_S} \cdot \frac{1}{4\pi p_c^3} x_o^{-2} \delta(x-x_o) \quad (8),$$

where p_o is some characteristic injection momentum and N_o/τ_S is the rate at which electrons are supplied to the acceleration region. The steady-state solution of equation (4) then is (Bogdan and Schlickeiser 1985)

$$f_\infty(p) = \frac{N_o}{16\pi p_c^3 \tau_S (\alpha D)^{1/2}} \frac{2^{a/2}\Gamma(b/2)}{\Gamma(2-(a/2))} x_o^{-2} x^{-a} \exp(-x^2/2)$$

$$\cdot \begin{cases} U(b/2, 2-(a/2), x_o^2/2)\, M(b/2, 2-(a/2), x^2/2) & \text{for } x \leq x_o \\ M(b/2, 2-(a/2), x_o^2/2)\, U(b/2, 2-(a/2), x^2/2) & \text{for } x \geq x_o \end{cases} \quad (9)$$

in terms of confluent hypergeometric functions.

Figure 1 illustrates the behaviour of $p^2 f_\infty(p)$ in different parameter regions deduced from the asymptotic forms of the confluent hypergeometric functions. At the very lowest ($p < p_o$) momentum [$p^2 f_\infty(p) \propto p^{2-a}$] and the very highest ($p \gg p_c$) momentum [$p^2 f_\infty(p) \propto p^{2-a-b} \exp(-p^2/(2p_c^2))$] the spectrum is rather independent of the values of a and b. At intermediate momenta, the spectrum is determined by b, the ratio of radiation loss time (τ_R) to escape time (T). If $b \leq 1$, i.e. $\tau_R < T$, the p^{2-a} power law extends up to the injection momentum p_o, and a constant spectrum ($a \leq 2$)

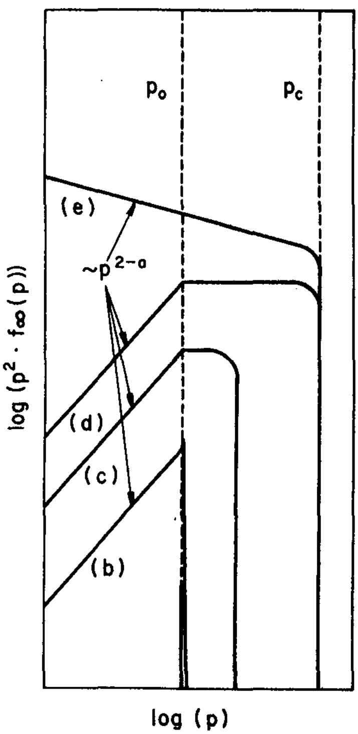

Fig. 1: Schematic illustration of some representative steady-state electron spectra $p^2 f_\infty(p)$ from (9). The relative normalizations of the spectra are arbitrary. The injection momentum (p_o), and the critical momentum p_c [Eq. (5)] are indicated by dotted lines. For spectrum (b) $b \gtrsim (p_c/p_o)^2$, $a < 2$; (c) $1 \ll b \ll (p_c/p_o)^2$, $a < 2$; (d) $b \lesssim 1$, $a < 2$; and (e) $b \lesssim 1$, $2 < a < 4$.

or p^{2-a} $(2 \leq a < 4)$ power law extends up from the injection momentum to the critical momentum p_c. In this case, the escape time is large and the spectrum is determined by the interplay between acceleration and radiation losses. If $b \lesssim 1$, i.e. $T < \tau_R$, the spectrum develops an $\exp(-b^{1/2}p/p_c)$ shoulder before the $\exp(-p^2/(2p_c^2))$ fall off for $p < p_c$. Finally when $b \gg 1$, and the escape time decreases still further, the spectrum develops a pronounced spike at the injection momentum. In this case, the particles escape from the acceleration region so quickly that they have little chance to interact with the magnetized plasma. Summarizing: electrons can be accelerated efficiently in flare stars if in the acceleration region the parameter b (equation (7)) is smaller than unity. The maximum momentum then is given by p_c in equation (5) which is that momentum where the Fermi acceleration time ($\tau_F = p/D$) exactly equals the radiation loss time ($\tau_R = 1/(\alpha p)$).

Another important point to note is, that the upper cutoff in the momentum spectrum (9) is due to synchrotron losses only in cases where $b \leq 1$. In the opposite situation $b \gtrsim 1$, the escape of energetic electrons from the acceleration region causes the cutoff. This should be kept in mind when interpreting the peak microwave emission of flares from RS CVn stars. Misidentifying the physical mechanism for the cutoff may yield wrong constraints on the physical conditions in the flare site, n, ℓ, λ and B.

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