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STOCHASTIC ACCELERATION OF SOLAR FLARE PROTONS

by

D.D. Barbosa

National Aeronautics and Space Administration
NGL 05-020-272

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ABSTRACT

We consider the acceleration of solar flare protons by cyclotron damping of intense Alfven wave turbulence in a magnetic trap. The energy diffusion coefficient D_{EE} is computed for a near-isotropic distribution of super-Alfvenic protons and a steady-state solution for the particle spectrum is found for both transit-time and diffusive losses out of the ends of the trap. The acceleration time to a characteristic energy $\sim$ 20 Mev/nucl can be as short as 10 sec. On the basis of phenomenological arguments we infer that the Alfven wave spectrum has a ω^{-2} frequency dependence and that the correlation time of the turbulence lies in the range 5×10^{-4}s $< \tau_{corr} < 5 \times 10^{-2}$ s.

I. INTRODUCTION

Solar flare protons are believed to be accelerated in complex magnetic bipolar configurations in active regions. Evidence points to the idea that acceleration occurs in a second stage process [De Jager, 1969] in contrast to primary energy release of the flare in a first stage process which is manifest directly in quasi-thermal heating and non-thermal 10-500 Kev electrons [Svestka, 1976]. The association of Type II radio bursts with proton flares is compatible with the scenario that a noisy MHD shock is generated soon ($\lesssim$ 3 mins) after the initiation of the flash phase and accelerates protons by stochastic processes as it propagates upwards and away from the flare site [McLean et al., 1971; Lin and Hudson, 1976]. The shock is certainly energetic enough and current belief is that Fermi processes are at work in conjuction with wave-particle scattering to energize protons up to energies > 10 Gev [Melrose, 1974; Svestka, 1976]. A large part of our theoretical understanding of acceleration mechanisms has derived from efforts directed towards understanding the energization of galactic cosmic rays [Rosen, 1969].

In this paper we will attempt to model the solar proton acceleration process with a very simple field-plasma configuration which might be described as a "leaky trap". Energization takes place in the corona in closed magnetic field structures but allowance is made for particle losses. As a first step in this direction we have considered only particle escape along the flux tube to the low solar atmosphere. The acceleration mechanism is cyclotron damping by super-Alfvenic protons

of intense Alfven waves propagating in the trap. This process is intrinsically stochastic and no adiabatic, Fermi-type mechanism is invoked for the acceleration.

Davis [1956] gave the earliest description of the effect that a turbulent state of MHD waves produces a diffusion in pitch angles towards isotropy. The development of quasi-linear theory [see Kennel and Englemann, 1966, and their refs.] provided a useful analytic tool with which one could address the problem of diffusion in phase-space resulting from turbulent wave-particle interactions. Sturrock [1966], using a Fokker-Planck formalism, achieved similar results and thus fortified our confidence in these methods.

Melrose [1974] demonstrated that the combined effects of stochastic pitch angle scattering by Alfven waves and Fermi acceleration could account for the energization of solar flare protons. The source of the Alfven turbulence was considered to be the non-thermal protons themselves which, if sufficiently anisotropic in pitch angle distribution, would drive an electromagnetic Alfven wave instability. Since the free energy for the waves is derived from the kinetic energy of the non-thermal protons in this picture, another, independent, free energy source must energize the protons (e.g., a Fermi process). In our model the energy flow is opposite to that of Melrose. We assume that Alfven turbulence is generated ad hoc by an exciting agency that is flare related. Using this approach we may eliminate the Fermi mechanism and consider the acceleration possible from just short wavelength, $kv \approx \Omega_{cp}$,

Alfven waves which not only scatter in pitch angle but also energize through collisionless damping [Schatzman, 1967; Sturrock, 1974].

The model can be applied to the system described by, for instance, McLean et al. [1971] in which a localized shock generates Alfven waves propagating throughout a flux tube. With some modifications, the model may also be applied to a system like that of Sturrock [1974] in which protons are accelerated on open field lines by a noisy MHD shock. However, recent γ-ray observations have suggested that some protons may be accelerated as early as the impulsive phase. We have, therefore, concentrated on the consequences of a leaky trap without pinning down the source of the turbulence or when and how it is generated. If impulsive phase proton acceleration is confirmed, then we hope our model may provide insight into the primary energy release mechanism itself.

In Section II we outline the ingredients of the model and inquire into the possibility of proton acceleration during primary energy release. In Section III we derive the energy diffusion coefficient, D_{EE}, for the cyclotron damping mechanism for a near-isotropic distribution of super-Alfvenic ions. In Section IV we apply this result to the leaky trap with either ballistic transit-time losses or diffusive losses out the ends of the trap. In the final section we summarize our results and comment on the possibility of flare-generated MHD standing waves in the trap.

II. MODEL

A. Basic Assumptions

The basic magnetic field configuration we shall consider is that of a closed loop, bipolar arch structure. We presume that some exciting agency, that is flare related, generates intense MHD turbulence which interacts with particles at coronal heights. The closed field lines serve as a magnetic trap for the non-thermal particles which are stochastically accelerated by the turbulence but are prevented from escaping by mirror forces arising from the convergence of the mean field $\vec{B}_0$ towards the photosphere. The turbulence is assumed to consist of short wavelength perturbations, $\lambda << L$, much smaller than any macroscopic scale of the system; this enables us to use homogeneous quasi-linear theory to gauge the acceleration efficiency. The dominant interaction with the waves is considered to be with the wavevector component parallel to $\vec{B}_0$ [i.e., $k_\perp \simeq 0$] and we, therefore, model the turbulence as an ensemble of parallel-propagating Alfven waves in the frequency range $\omega_0 < \omega << \Omega_{cp}$, where Ω_{cp} is the proton gyrofrequency and ω_0 is a low frequency cutoff.

The question of particle phase-space losses is an important one and different assumptions concerning what loss mechanism prevails will alter the picture significantly. We shall consider proton energies sufficiently large that they may be considered collisionless with respect to binary coulomb interactions [Braginskii, 1965] during the acceleration period. However, if the protons remain in dense regions sufficiently long, this deenergization process must be accounted for. We may expect that some cross-field diffusion of particles out of the trap is likely to occur during the acceleration period. However, we have chosen to consider what

may be the fastest loss process: particle escape along the field to the chromosphere and photosphere. A consequence of our assumption of intense Alfven turbulence (with a predominant magnetic component) is that particles will be scattered very efficiently in pitch angle and can either stream out or diffuse out of the acceleration region along the flux tube to the lower atmosphere. For either case we have scaled the losstime to the particle ballistic travel time and the diffusion time along the length of the flux tube. In the first case (transit-time losses) our model is analogous to the process described by Kennel [1969] who considered stochastic acceleration of electrons in the earth's magnetosphere. In the second case (diffusive losses) the model incorporates diffusive spatial transport in the presence of elastic wave collisions [Jokipii, 1971] analogous to the propagation effect observed in interplanetary space.

B. Applications

The association of Type II radio bursts with proton events observed at earth suggests that a flare-initiated shock propagates upwards and accelerates protons. If the shock is the exciting agent of MHD waves, then our model may be used to compute the acceleration of protons in a configuration like that proposed by McLean et al. [1971]. It must be determined, for instance, if the acceleration takes place ahead of or behind the shock, and the question of how the protons obtain access to interplanetary field lines must be addressed. If the field lines of the trap "open up" [Kopp and Pneuman, 1976], then particles will be injected ballistically into the solar wind; but if the field lines remain closed, but distended, cross-field diffusion effects must occur to exhaust the

trap. Chronologically, this process is an "after effect" of the flare; that is, subsequent to or during the primary energy release the shock is formed and proton acceleration begins at times $\gtrsim$ 3 mins after the initiation of the flash phase.

However, γ-ray observations [Chupp et al., 1973] made by the OSO-7 satellite have indicated that proton acceleration may begin at times coincident with the flash phase onset. Figure 1 is a plot of the Aug. 4, 1972 2.2 Mev γ-ray line together with hard x-ray profiles from Van Beck et al. [1973]. We note that while the maximum of the γ-ray flux occurs after the hard x-ray maximum a not insignificant flux is measured very soon after onset [Svestka, 1976]. In fact, the flux at 0625 UT is 1/3 of the maximum. Since there must be a finite acceleration time for the protons to reach these energies in sufficient number, Figure 1 indicates that, within the temporal resolution of the instrument and ignoring any nuclear reaction time delay, a generous upper limit for the acceleration time $\sim$ 3 mins (assuming that energization begins at 0622). The acceleration time could be much less (10s $< \tau_{acc} <$ 200s) and still be consistent with the non-thermal 2.2 Mev emission measure reaching a maximum later as more flux tubes become involved and the acceleration region increases. The measurement at 0622, if real, gives support for a short $\sim$10s acceleration time. Ramaty et al. [1975] interpret a break in the hard x-ray spectrum of Aug. 4, 1972 at about 500 Kev towards a harder profile at higher energies as evidence for electrons undergoing second stage acceleration.

Since we might expect intense Alfven wave turbulence generated in primary energy release mechanisms like, for instance, magnetic field tearing instabilities [Spicer, 1976], we have, therefore, considered a parameter regime for the leaky trap that might be applied to plasma-field conditions in the vicinity of x-ray kernals and their coronal loops at the time of the impulsive phase. The model, then, can allow us to infer something about the primary energy release process itself rather than post-explosion conditions in MHD shocks. Future observations should define clearly the role of magnetic turbulence, especially with regard to primary energy release, and can test our supposition that proton acceleration can occur as early as the impulsive phase.

III. DIFFUSION COEFFICIENT

A. Quasi-linear Theory

We consider a homogeneous plasma in a magnetic field $\vec{B}_o$ directed along the z axis and assume that the turbulence, composed of parallel-propagating Alfven waves ($k_\perp = o$) with frequencies in the range

$$\omega_o < \omega = |k| V_A << \Omega_{cp}, \tag{1}$$

can be described by a spectral density

$$< B_1^2/8\pi > = 2 \int_o^\infty d\omega \, \mathcal{B}(\omega). \tag{2}$$

$V_A = B_o/(4\pi M n_o)^{\frac{1}{2}}$ is the Alfven velocity, $\Omega_{cp} = eB_o/Mc$ is the proton gyrofrequency, and $\vec{B}_1$ is the perturbation magnetic field related to the induced electric field by

$$|\vec{E}_1| = \frac{V_A}{c} |\vec{B}_1|. \tag{3}$$

The factor of two in equation (2) is to discriminate between waves travelling parallel and antiparallel to $\vec{B}_o$.

It is well known that the temporal evolution of a distribution of particles, say protons, in response to resonant scattering by Alfven waves, can be described by a quasi-linear relation [Kennel and Engleman, 1966; Hall and Sturrock, 1967; Davidson, 1972] which, for $k_\perp = o$, reduces to in the non-relativistic regime

$$\frac{\partial}{\partial t} f(\vec{v},t) = -\frac{\partial}{\partial \vec{v}} \cdot \vec{J} = -\frac{\partial J_z}{\partial v_z} - \frac{1}{v_\perp} \frac{\partial}{\partial v_\perp} (v_\perp J_\perp) \tag{4}$$

where

$$\vec{J} = -\frac{4\pi^2 e^2 V_A^2}{M^2 c^2} \frac{|V_g| \mathcal{B}(\omega)}{|V_g - v_z|} \left[(1 - \frac{kv_z}{\omega}) \frac{\partial f}{\partial v_\perp} + \frac{kv_\perp}{\omega} \frac{\partial f}{\partial v_z} \right] \left[(1 - \frac{kv_z}{\omega}) \hat{v}_\perp + \frac{kv_\perp}{\omega} \hat{v}_z \right]. \tag{5}$$

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In the frequency range considered, the group velocity is the Alfven speed, $|V_g| = V_A$; $\hat{v}_z$ and $\hat{v}_\perp$ are unit vectors in velocity space (cylindrical coordinates), azimuthal symmetry about $\vec{B}_0$ being assumed. The effect of the finite perturbation electric field (3) is represented in the terms consisting of the integer 1. The unique relationship between the proton and the Alfven wave in this interaction is expressed in the cyclotron resonance condition which is given as

$$\omega - kv_z = \pm\Omega_{cp} \ , \tag{6}$$

the upper (lower) sign taken for the L (R) mode of the Alfven wave. For $k > o$ the terms L and R also give the polarization of the wave. Equation (6) indicates that an effective "collision" occurs when the doppler-shifted wave frequency is the proton's gyrofrequency and the sense of rotation is the same.

For the range of parameters we shall consider, $V_A \ll c$ and the effect of the predominantly magnetic wave is to scatter the particles primarily in pitch angle. It is convenient, therefore, to express (5) in spherical coordinates. If $\cos\theta = v_z/v$ we have, using (6),

$$\vec{J} = -\frac{4\pi^2 e^2 V_A^2}{M^2c^2}\frac{\Omega_{cp}^2}{\omega^2}\frac{|V_g|\mathcal{B}(\omega)}{|V_g - v_z|}\left[\left(\sin\theta \pm \frac{kv_\perp}{\Omega_{cp}}\cos\theta\right)\frac{\partial f}{\partial v} + \left(\cos\theta \mp \frac{kv_\perp}{\Omega_{cp}}\sin\theta\right)\frac{1}{v}\frac{\partial f}{\partial\theta}\right]\left[\left(\sin\theta \pm \frac{kv_\perp}{\Omega_{cp}}\cos\theta\right)\hat{v} + \left(\cos\theta \mp \frac{kv_\perp}{\Omega_{cp}}\sin\theta\right)\hat{\theta}\right] . \tag{7}$$

Inspection of (6) indicates that if $\omega \ll \Omega_{cp}$, then $|v_z| \gg V_A$ if resonant scattering is to occur. Protons with $|v_z| < V_A$ follow the wave adiabatically conserving their first invariant $\mu = \frac{mv_\perp^2}{2B}$ during the

oscillation, whereas super-Alfvenic protons can satisfy (6) and are non-adiabatically scattered, the interaction being sufficiently rapid ($\tau \sim \Omega_{cp}^{-1}$) so as to violate their first invariant. Equation (6) then permits a useful perturbation expansion [Melrose, 1974]

$$k/\Omega_{cp} = \frac{\mp 1/v_z}{1-\omega/kv_z} = \mp \frac{1}{v_z}\left[1 + \frac{\omega}{kv_z} + \cdots\right] \tag{8}$$

for particles with velocities $|v_z| >> V_A$. Retention of the zero order term in (8) provides strictly magnetic effects, while the first order term incorporates effects of a finite electric field due to $\omega \neq o$. The expansion of (7) to zero order yields

$$\vec{J}^{(o)} = -\frac{4\pi^2 e^2}{M^2c^2} V_A \mathcal{B}(\omega) \frac{1}{|\cos\theta|} \frac{\partial f}{\partial \theta} \hat{\theta} \ . \tag{9}$$

Any reference to frequency here is incidental since if (2) is given alternately as

$$< B_1^2/8\pi > = \int_{-\infty}^{\infty} dk\, \tilde{\mathcal{B}}(k) \tag{10}$$

we have the relation

$$V_A \mathcal{B}(\omega) = \tilde{\mathcal{B}}(k) \tag{11}$$

and thus (9) can also represent non-adiabatic pitch angle scattering from stationary magnetic inhomogeneities described by $\tilde{\mathcal{B}}(k)$ [Jokipii, 1971].

On the shortest time scale, the turbulence acts to isotropize the distribution and on a longer time scale $O(v^2/V_A^2)$ the particles are energized through stochastic accelerations from $\vec{E}_1$. This results in a diffusion current $J_v^{(2)}$ which we may conveniently write using the approximation that the distribution is nearly isotropic from pitch angle scattering and $\frac{\partial f}{\partial \theta} \simeq o$. Retaining the first order term in (8) provides the second order term for $\vec{J}^{(2)} = O(V_A^2/v^2)$

$$\mathcal{J}^{(2)} = \frac{-4\pi^2 e^2 V_A^2}{M^2 c^2} \mathcal{B}(\omega) \frac{V_A \sin^2\theta}{v|\cos\theta|} \frac{\partial f}{\partial v} \hat{v} \tag{12}$$

and the distribution evolves as

$$\frac{\partial}{\partial t} f(v,t) = \frac{4\pi^2 e^2 V_A^3}{M^2 c^2} \frac{\sin^2\theta}{|\cos\theta|} \frac{1}{v^2} \frac{\partial}{\partial v}\left[\mathcal{B}(\omega)\, v \frac{\partial f}{\partial v}\right]. \tag{13}$$

B. Wave Spectrum

To model the turbulence we assume the spectral density can be represented by a power law in frequency

$$\mathcal{B}(\omega) = \mathcal{B}_0\, \omega^{-\delta}\, \theta(\omega-\omega_0) \ . \tag{14}$$

$\theta(x)$ is the Heaviside step function which introduces a low frequency cutoff at ω_0. It is also implicit in (14) that frequencies do not extend as high as the proton gyrofrequency. The resonance condition (6) gives the inverse wave-particle relation

$$\omega = |k|V_A \simeq \Omega_{cp} \frac{V_A}{v|\cos\theta|} \ . \tag{15}$$

The subsequent steps are straightforward: the spectral density is normalized to the perturbation field energy density in (2) to determine $\mathcal{B}_0$, equations (14) and (15) are inserted into (13) and the equation is averaged over pitch angles, and the change of variable to energy per nucl is made.

Generalizing this procedure for an arbitrary ion of ionic charge Ze and mass AM, the temporal evolution of the differential number density for any ionic species can be written as

$$\frac{\partial}{\partial t} N(E,t) = D \frac{\partial}{\partial E}\left[E^{n+3/2} \frac{\partial}{\partial E}\left(\frac{N}{\sqrt{E}} \right) \right] \tag{16}$$

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where the total number density of the ionic species is given as

$$N = \int_0^\infty dE\ N(E) \tag{17}$$

and

$$D = (A/Z)^{2n-1}\ D_p\ , \tag{18}$$

$$D_p = 2\pi\Omega_{cp}\ \frac{B_1^2/8\pi}{B_0^2/8\pi}\ \frac{2n}{(2n+1)(2n+3)}\ \frac{E_A}{E_{mp}^n} \tag{19}$$

The variable $E = \frac{1}{2}Mv^2$ = energy/nucl of the species, $E_A = \frac{1}{2}MV_A^2 = \frac{B_0^2}{8\pi n_0}$ is the Alfven energy, M = proton mass, and the index n is related to the wave spectral index by

$$n = \tfrac{1}{2}(\delta-1)\ . \tag{20}$$

The normalization of $\mathcal{A}(\omega)$ introduced the low frequency cutoff ω_0 into the scaling of the diffusion coefficient and we have written it in energy units as

$$E_{mp} \equiv \tfrac{1}{2}MV_A^2 \left(\frac{\Omega_{cp}}{\omega_0}\right)^2 \tag{21}$$

E_{mp} is related to the maximum energy a proton can achieve in the following way: if E_{mp} is non-relativistic in value, then $E_{max,p} = E_{mp}$, but if E_{mp} is relativistic, one must use the relativistic form of the resonance condition instead of (6) to determine $E_{max,p}$. Thus, the maximum energy a proton can be accelerated to is

$$E_{max,p} = \begin{cases} E_{mp} & : \ E_{mp} << 1\text{Gev} \\ Mc^2\ \dfrac{V_A}{c}\ \dfrac{\Omega_{cp}}{\omega_0} & : \ E_{mp} >> 1\text{Gev} \end{cases} \tag{22}$$

Again, so long as we are considering the non-relativistic regime of energies in (16), (21) is used by definition for all values of ω_0.

It is of interest to ask what minimum energy (injection energy) of electrons is necessary for them to be accelerated by Alfven waves. The

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resonance condition for electrons with gyrofrequency $\Omega_{ce} = eB_o/mc$ is given as

$$\omega - kv_z = \mp \Omega_{ce}/\gamma \tag{23}$$

which indicates that if $\omega < \Omega_{cp}$, we have [Sturrock, 1974]

$$E_{inj,p} = E_A = \tfrac{1}{2}MV_A^2$$

$$E_{inj,e} = \begin{cases} \frac{M}{m} E_A & : \text{non-rel} \\ Mc^2 \frac{V_A}{c} & : \text{ultra-rel} \end{cases} \tag{24}$$

These restrictions apply for MHD Alfven waves where $\omega < \Omega_{cp}$. Extension of the wave spectrum to higher frequencies to include whistler waves (R mode) or Langmuir waves lowers the electron energy threshold [Melrose, 1974], but is better addressed as a separate problem since the wave properties are quite different and, phenomenologically, acceleration by higher frequency waves may be related to first stage acceleration processes [Hoyng, 1977].

If, instead of energy, one considers a diffusion in the non-relativistic regime of rigidity, $R = AMvc/Ze$, we have alternatively

$$\frac{\partial}{\partial t} \tilde{N}(R,t) = K \frac{\partial}{\partial R} \left[R^{2n+2} \frac{\partial}{\partial R} (\tilde{N}/R^2) \right] \tag{25}$$

where

$$K = \frac{A}{Z} \frac{\pi}{2} \Omega_{cp} \frac{B_1^2/8\pi}{B_o^2/8\pi} \frac{2n}{(2n+1)(2n+3)} \frac{R_A^2}{R_{mp}^{2n}} \quad . \tag{26}$$

C. Physical Interpretation

An interpretation of the physical process leading to equations (4), (6), and (14) is given in Figure 2. We have decomposed the turbulence into an ensemble of elemental scattering centers of length $x \leq x_{corr}$. This ensemble is equivalent to the actual situation in that the idealized

scatterers have the same autocorrelation function as the actual system,

$$\rho(t) = \int_{-\infty}^{\infty} d\omega\, \mathcal{B}(\omega) e^{-i\omega t} / [\int_{-\infty}^{\infty} d\omega\, \mathcal{B}(\omega)] \tag{27}$$

$$\tilde{\rho}(x) = \int_{-\infty}^{\infty} dk\, \tilde{\mathcal{B}}(k) e^{ikx} / [\int_{-\infty}^{\infty} dk\, \tilde{\mathcal{B}}(k)] \quad . \tag{28}$$

From the power spectrum given in (14) we identify ($\delta > 1$)

$$\tau_{corr} \simeq \frac{1}{\omega_0} \qquad x_{corr} \simeq V_A\, \tau_{corr} \,, \tag{29}$$

and $\tilde{\mathcal{B}}(k)$ [or $\mathcal{B}(\omega)$] is a measure of the differential number density of elemental scattering centers of size $k^{-1} \sim x \leq x_{corr}$. A super-Alfvenic proton satisfying the resonance condition (6) can be effectively scattered as drawn in Figure 2. Particle A sees a component of $\vec{E}_1$ parallel to $\vec{v}$ throughout the interaction and is accelerated, while particle B is 180° out of phase and is decelerated. A collection of random (cyclotron) phased particles will experience, however, a dispersion in energies.

The acceleration process described here cannot be interpreted as a stochastic Fermic type. The acceleration mechanism described by Fermi [1949, 1954] involves an adiabatic process such that the particles' first and second invariants, $\mu = mv_\perp^2/2B$, $J = \oint v_{\parallel} ds$, are conserved. Pitch angle scattering is an ancillary process invoked to counterbalance the systematic decrease in the particle's pitch angle as it is accelerated [Davis, 1956]. We have described a process where oscillations are sufficiently rapid that they non-adiabatically violate the first invariant. Energization (or denergization) is concomitant with pitch angle scattering when $\omega \neq o$ and such a process is best interpreted as cyclotron damping of Alfven waves [Stix, 1962].

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IV. CHARACTERISTIC SCALINGS

In order to apply the result of equation (16) to the leaky trap model, we have considered three problems: a transient solution and a steady-state solution with transit-time and diffusive losses.

A. Initial Value Problem

Two facts which immediately emerge from injection of (16) is that the energization time scales as $\tau_E \sim E^{1-n}$ and that the solution asymptotes to a "steady state" $N(E) \sim E^{-n}$. The energization time indicates that $n = 1$ distinguishes two, quite different, diffusion regimes which favor acceleration of either high or low energy particles. We shall argue in this paper that the regime which favors low energy particle acceleration is more appropriate for solar proton flares and we shall consider the parameter regime

$$1 < \delta < 3 \leftrightarrow 0 < n < 1 . \tag{30}$$

This statement is equivalent to the expectation that if the power spectrum in (14) is sufficiently flat, protons can be accelerated over a large energy range in the shortest time for a given perturbation field energy density. Increasing δ to larger values puts all the energy in the low frequencies and biases the acceleration in favor of very high energy particles [Equation (15)].

The homogeneous solution of (16) with the boundary conditions at an injection energy E_0 that

$$N(E_0,t) = N_0 \ , \ N(E>E_0,t=0) = 0, \ N(E \to \infty) \to 0$$

is given as

$$N(E,t) = \frac{N_o}{2\pi i} \int_{-i\infty}^{i\infty} \frac{ds}{s} e^s \left(\frac{E}{E_o} \right)^p \frac{K_\mu(\xi E^w \sqrt{s})}{K_\mu(\xi E_o^w \sqrt{s})} , \tag{31}$$

with

$$p = \tfrac{1}{4}(1-2n) \qquad \xi = \frac{2}{(1-n)} \sqrt{1/Dt}$$
$$w = \tfrac{1}{2}(1-n) > 0 \qquad \mu = \frac{2n+1}{2(1-n)} \quad . \tag{32}$$

If we consider times long enough such that $\xi E_o^w << 1$, the small-argument form of $K_\mu(z)$ can be used in the denominator and the integral is evaluated as

$$N(E,t) \simeq N_o \left(\frac{E_o}{E} \right)^n \frac{\Gamma(\mu, \tfrac{1}{4}\xi^2 E^{2w})}{\Gamma(\mu)} \tag{33}$$

where $\Gamma(a,x)$ is the incomplete gamma function [Abromowitz and Stegun, 1970]. This function exhibits two different behaviors depending on whether

$$x = \tfrac{1}{4}\xi^2 E^{2w} \gtrless 1 . \tag{34}$$

We may approximate the solution as

$$N(E,t) \simeq \begin{cases} N_o \left(\frac{E_o}{E} \right)^n & : \ x << 1 \\ N_o \left(\frac{E_o}{E} \right)^n \frac{(\tfrac{1}{4}\xi^2 E^{2w})^{\mu-1}}{\Gamma(\mu)} e^{-\tfrac{1}{4}\xi^2 E^{2w}} & : \ x >> 1 \end{cases} \tag{35}$$

The solution changes from a power low at low energies to an exponential behavior (w>0) at the diffusion front characterized by a break energy, E_B, which depends on time and is defined as x = 1 leading to

$$E_B(t) = [(1-n)^2 \, Dt]^{\frac{1}{1-n}} \tag{36}$$

In a subsequent section we shall argue for the feasibility of the value $n = \frac{1}{2}$. Anticipating this discussion and defining $\epsilon_B = B_1^2/B_0^2$ we may scale the break energy as

$$E_B(n=\tfrac{1}{2},t) = 0.61 \left(\frac{B_o}{100G} \right)^6 \left(\frac{\epsilon_B}{10^{-2}} \right)^2 \left(\frac{10Gev}{E_{mp}} \right) \left(\frac{6x10^9 cm^{-3}}{n_o} \right)^2$$

$$t^2 \text{ Mev/nucl} . \tag{37}$$

Time is in units of seconds and our choice of B_o and n_o = background density corresponds to an Alfven energy of $E_A \approx 41.4$ Kev. Note also [Equation 18] that for $n = \frac{1}{2}$, reference to Z/A disappears and (37) is valid for any ionic species. The value 10 Gev for E_{mp} corresponds to $E_{max\ p} \approx$ 4.3 Gev.

B. Steady-State Solution

Assuming that proton acceleration can occur for a period longer than 10 sec, (37) admits the possibility of very efficient proton acceleration, E_B being a measure of the "hardness" of the proton spectrum. A realistic model should allow for particle losses and phenomenologically we may write

$$\frac{\partial N}{\partial t} = D \frac{\partial}{\partial E} \left[E^{n+3/2} \frac{\partial}{\partial E} \left(\frac{N}{\sqrt{E}} \right) \right] - \frac{N}{\tau_o E^\beta} = 0 \tag{38}$$

to describe the time asymptotic evolution of the transient solution (33). The steady-state solution of (38) should give a meaningful result provided that the fastest losstime

$$\tau = \tau_o E^\beta \tag{39}$$

can be ascertained. The solution for $q = \frac{1}{2}(1-n-\beta) > 0$ is

$$N(E) = N_o \left(\frac{E}{E_o} \right)^p \frac{K_\nu(\lambda E^q)}{K_\nu(\lambda E_o^q)} \quad , \tag{40}$$

where

$$p = \tfrac{1}{4}(1-2n) \qquad \lambda = \frac{2}{(1-n-\beta)} \sqrt{\frac{1}{D\tau_o}}$$

$$q = \tfrac{1}{2}(1-n-\beta) > 0 \qquad \nu = \frac{2n+1}{2(1-n-\beta)} \quad . \tag{41}$$

If $\lambda E_o^q << 1$ we may use small argument form for the denominator and (40) exhibits the following behavior

$$N(E) \simeq \begin{cases} N_o \, (E_o/E)^n & : \lambda E^q << 1 \\ N_o \dfrac{\sqrt{\pi}}{\Gamma(\nu)} \left(\tfrac{1}{2}\lambda E_o^q \right)^{\nu - \frac{1}{2}} \left(\dfrac{E_o}{E} \right)^{\frac{1}{4}(n-\beta)} e^{-\lambda E^q} & : \lambda E^q >> 1 \quad . \end{cases} \tag{42}$$

The steady-state break energy, defined as $\lambda E^q = 1$, is

$$E_B = [\tfrac{1}{4}(1-n-\beta)^2 \, D\tau_o]^{\frac{1}{1-n-\beta}} \quad . \tag{43}$$

C. Transit-Time Losses

In a magnetic field minor geometry, particles with pitch angles less than the loss cone value

$$\sin^2\theta_o = B_o/B_{max} \tag{44}$$

can penetrate to a field strength of B_{max} where they may be considered as lost [through significant denergization processes in the dense chromosphere and photosphere]. This angle corresponds to a loss cone "volume" of $4\pi[1-\cos\theta_o]$ steradians. If pitch angle scattering occurs rapidly enough to maintain a significant number of particles in the loss cone, then the

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losstime is estimated to be the transit time along the flux tube with a weighting factor of the fractional volume of the loss cone [the strong pitch angle diffusion regime of Kennel, 1969], i.e.,

$$\tau = (M/2)^{\frac{1}{2}} L(1-\cos\theta_0)^{-1} E^{-\frac{1}{2}} \quad . \tag{45}$$

Transit-time losses lead to a value of $\beta = -\frac{1}{2}$ in (39). We note that if $n = \frac{1}{2}$, then the solution (40) behaves at large energies as $N(n=-\beta=\frac{1}{2}) \sim E^{-\frac{1}{4}} e^{-\lambda\sqrt{E}}$. Observations made at 1AU have indicated that over a large range of energies, proton spectra can be fit by an exponential in rigidity, the fit improving at high energies ($\gtrsim$200MV $\simeq$ 21Mev) [Freier and Webber, 1963; Van Hollebeke et al., 1975]. An exponential in rigidity is also consistent with models based on γ-ray observations [Ramaty and Lingenfelter, 1972]. On this basis we argue for the feasibility of the value $n = \frac{1}{2}$ corresponding to a wave spectral index of $\delta = 2$.

If we take a coronal field strength of 100G and a photospheric field of 1000G, the loss cone angle $\theta_0 \simeq 18°$. Using (45) in (43) gives a steady-state break energy for transit-time losses of

$$E_B\ (n=-\beta=\tfrac{1}{2}) = 11 \left(\frac{B_0}{100G}\right)^3 \left(\frac{\varepsilon_B}{10^{-2}}\right) \left(\frac{L}{10^9 cm}\right) \left(\frac{10Gev}{E_{mp}}\right)^{\frac{1}{2}} \left(\frac{6x10^9 cm^{-3}}{n_0}\right)$$

$$\left(\frac{0.1}{B_0/B_{max}}\right) \text{Mev/nucl} \quad . \tag{46}$$

Although (46) gives a reasonable value for the turnover, the scaling may not be correct. In the absence of cross-field losses, (45) without the loss cone weighting factor must be the shortest conceivable time to deplete the trap. But, for the parameters chosen, the mean free path for

pitch angle scattering is much less than the length of the tube. Under these circumstances, a spatial diffusion along the field to the lower atmosphere may be more appropriate than losses scaled to the particle ballistic transit time.

D. Diffusive Losses

If the phase-space distribution function varies along the z direction and we consider only the effects of pitch angle scattering by waves, then (4) in zero order can be written as

$$\frac{\partial f}{\partial t} + v_z \frac{\partial}{\partial z} f\,(z,\vec{v},t) = \frac{\partial}{\partial \eta}\left[\, D_{\eta\eta} \frac{\partial f}{\partial \eta} \right] \tag{47}$$

and $\eta = \cos\theta$. Jokipii [1971] has shown that in the approximation that f is nearly isotropic, (47) reduces to a diffusion equation for the differential number density

$$\frac{\partial}{\partial t} N(z,E,t) = K_{\parallel} \frac{\partial^2 N}{\partial z^2} \quad , \tag{48}$$

where

$$K_{\parallel} = \frac{2}{9} v^2 \left[\int_{-1}^{1} d\eta\; D_{\eta\eta} \right]^{-1} \tag{49}$$

and from (9) we have for protons

$$D_{\eta\eta} = \frac{4\pi^2 e^2}{M^2 c^2} \mathcal{B}(\omega) \frac{V_A}{v} \frac{1-\eta^2}{|\eta|} \quad . \tag{50}$$

Substituting the spectral density (14) in the above and generalizing for an arbitrary ion yields the result

$$\int_{-1}^{1} d\eta\; D_{\eta\eta} = \overline{D}_{\eta\eta} = \left(\frac{A}{Z}\right)^{2n-1} \pi\Omega_{cp} \frac{B_1^2/8\pi}{B_o^2/8\pi} \frac{2n}{(2n+1)(2n+3)} \frac{E^n}{E_{mp}^n} \quad , \tag{51}$$

and the parallel spatial diffusion coefficient is

$$K_{\parallel} = \frac{4}{9} \frac{E}{M \overline{D}_{nn}} \quad . \tag{52}$$

If we consider ions diffusing down both legs of a trap of length L and estimate from (48) the losstime as $\tau_{diff} \simeq \frac{L^2}{4K_{\parallel}}$, we have

$$\tau_{diff} \simeq \frac{9}{16} ML^2 \left(\frac{A}{Z} \right)^{2n-1} \pi\Omega_{cp} \, \epsilon_B \frac{2n}{(2n+1)(2n+3)} \frac{E^{n-1}}{E_{mp}^{n}} \quad . \tag{53}$$

If $n = \frac{1}{2}$ then $\beta = -\frac{1}{2}$ and the steady state solution will behave in a manner similar to the results of the previous section for transit time losses. Equations (53) and (43) give a turnover for diffusive losses of

$$E_B \; (n=\tfrac{1}{2}) = 17.2 \left(\frac{B_o}{100G} \right)^4 \left(\frac{\epsilon_B}{10^{-2}} \right)^2 \left(\frac{L}{10^9} \right)^2 \left(\frac{6\times10^9}{n_o} \right) \left(\frac{10Gev}{E_{mp}} \right)$$

Mev/nucl . (54)

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V. SUMMARY AND DISCUSSION

We have considered some of the consequences of a turbulent spectrum of intense Alfven waves propagating along the magnetic field of coronal flux tubes. The model describes a steady-state situation where a population of ions is maintained at a threshold velocity of the Alfven speed and is scattered in pitch angles by the turbulence towards isotropy and energized concomitantly through collisionless damping. The acceleration time to $\sim$10Mev energies can be as short as $\sim$10s and the process itself is limited only by how well the trap can contain the accelerated particles. The measure of the acceleration efficiency has been gauged by a break energy of the particle spectrum which characterizes a turnover from a relatively flat partical spectrum to a rapidly decreasing exponential at high energies. This break energy is approximately the value at which the acceleration time equals the loss time.

The primary loss mechanisms which we have considered are transit-time losses and diffusive losses out of the ends of the flux tube. For the range of parameters such that the mean-free-path for pitch angle scattering is small relative to the length of the tube, the diffusive picture is probably more correct. The question of other losses is important and further work is needed to clarify the role of cross-field diffusive effects, especially with regard to how particles obtain access to interplanetary field lines. However, if the acceleration region consists of open field lines [Sturrock, 1974] this problem disappears.

One important conclusion we have drawn is that for a given r.m.s. magnetic perturbation B_1 the wave power spectrum in (14) must be relatively

flat (30) if efficient acceleration is to occur. For large values of the wave spectral index δ the diffusion process favors the acceleration of the very high energy particles which are scattered by the most energetic oscillations near the correlation length of the turbulence; in this case the energy diffusion coefficient scales as a strong power of $(E_0/E_{mp})^n$ [Newman, 1975] indicating that the acceleration of low energy particles near the injection energy E_0 is inefficient and holding up the works. On the basis of this model proton flares are associated with both large magnetic field perturbations and flat power spectra (30). This result is consistent with the findings of La Combe [1977] who considered particle acceleration in radio galaxies by Alfven waves. We may also extend the result (19) to include the value $\delta = 1 (n=0)$ for which the substitution is made in both (19) and (51)

$$\frac{2n}{(2n+1)(2n+3)} \xrightarrow[\delta = 1]{} \frac{1}{3} \quad \frac{1}{\ln \omega_{max}/\omega_0} \tag{55}$$

In particular, we have shown that the value $\delta = 2$ leads to an exponential particle spectrum at high energies that behaves as $N \sim E^{-\frac{1}{4}} e^{-\lambda\sqrt{E}}$. On the basis of observational indications of this behavior we have assumed $\delta = 2$ as a likely value. For this case, we may also rewrite the results (52) and (54) in a simple form in terms of the low frequency cutoff ω_0 and we offer the following speculations on second stage acceleration in conclusion:

The choice of scaling the diffusion coefficients in terms of the quantity E_{mp} (21) throughout this paper was motivated by the possibility of determining the correlation time, ω_0^{-1}, from observations of the highest

recorded proton energies. The value taken of E_{mp} = 10Gev corresponds to $E_{max,p} \approx 4.3$ Gev and from (21) we infer a correlation time and distance of

$$\tau_{corr} = 5.1\times10^{-4} \left(\frac{100G}{B_o}\right)^2 \left(\frac{n_o}{6\times10^9}\right)^{\frac{1}{2}} \left(\frac{E_{mp}}{10Gev}\right)^{\frac{1}{2}} \text{ sec}$$

$$x_{corr} = 1.5 \left(\frac{100G}{B_o}\right) \left(\frac{E_{mp}}{10Gev}\right)^{\frac{1}{2}} \text{ km} \quad . \tag{56}$$

It is possible, however, that the correlation time is much longer than (56) and the much higher proton energies possible are not recorded because of the rapid decrease in the spectrum (40) at very high energies. Treating ω_o then as a free parameter and taking n = ½, we may rewrite in summary the results for spatial diffusion as $[\varepsilon_B = B_1^2/B_o^2]$

$$D_{\eta\eta}(n=\tfrac{1}{2}) = \frac{\pi}{8}\,\omega_o\,\varepsilon_B\,\frac{v}{V_A} \tag{57}$$

and

$$K_{||}(n=\tfrac{1}{2}) = \frac{16}{9\pi}\,\frac{v}{k_o\varepsilon_B} \quad . \tag{58}$$

Defining a mean-free-path such that $K_{||} = \frac{1}{3}\,v\,\lambda$ we have for spatial diffusion

$$\lambda(n=\tfrac{1}{2}) = \frac{16}{3\pi}\,\frac{1}{k_o\varepsilon_B} \quad . \tag{59}$$

The energy diffusion coefficient (19) becomes

$$D(n=\tfrac{1}{2}) = \frac{\pi}{4}\,\omega_o\,E_B\,E_A^{\frac{1}{2}} \tag{60}$$

leading to the temporal break energy (36) of

$$E_B(n=\tfrac{1}{2},t) = \left(\frac{\pi}{16}\right)^2 \varepsilon_B^2\,E_A\,(t/\tau_{corr})^2 \tag{61}$$

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and a steady-state break energy for diffusive losses of

$$E_B(n=\tfrac{1}{2}) = \frac{9\pi^2}{2048}\,\omega_o^2\,\epsilon_B^2\,ML^2 \quad . \tag{62}$$

If we estimate from (16) an energy diffusion rate

$$\frac{1}{\tau_{EE}} \sim \frac{D}{E^{\frac{1}{2}}} = \frac{\pi}{4}\,\omega_o\,\epsilon_B\,\frac{V_A}{v} \tag{63}$$

we get a result which differs from (61) by a factor of 4 but is equal to Melrose's [1974] result [see his equation 24 and the discussion following]. Equation (62) may be scaled as

$$E_B(n=\tfrac{1}{2}) = 18\left(\frac{0.05s}{\tau_{corr}}\right)^2\left(\frac{\epsilon_B}{1}\right)^2\left(\frac{L}{10^9cm}\right)^2 \text{Mev/nucl} \quad , \tag{64}$$

and it is clear that longer correlation times than (56) require a larger ϵ_B in proportion to achieve a break energy in the range of 20Mev/nucl. In fact if we consider that the value (64) for the break energy is typical we may place the correlation time in the range

$$5\text{x}10^{-4}s < \tau_{corr} < 5\text{x}10^{-2}s \quad , \tag{65}$$

where the smaller limit is determined by the parameters given by (56) such that $E_{max,p} > 4.3$Gev and the larger limit is determined by (64) such that $\epsilon_B < 1$.

The limit $\epsilon_B = 1$ offers an interesting interpretation in terms of low frequency MHD standing waves of the trap. If the diffusive scale length is L/2 and the wavenumbers of the magnetic perturbations take the harmonic values $k = 2\pi n/\ell$ where $\ell < L$, the break energy for diffusive losses (62) can be written as

$$E_B(n=\tfrac{1}{2}) = \frac{9\pi^4}{256}\,\epsilon_B^2\left(\frac{L}{\ell}\right)^2 n_{corr}^2\,E_A \tag{66}$$

where $k_o = 2\pi n_{corr}/\ell$ is the dominant wavenumber excited. If we require an

energization of about $E_B = 10^3\ E_A$, and suppose that, say, $L = 3\ell$, then the limit $\epsilon_B = 1$ is consistent with a low harmonic, $n_{corr} \simeq 6$, being excited during the acceleration. The fundamental frequency, $\omega_{min} = 2\pi V_A/\ell$, gives a greatest upper bound to the correlation time of

$$\tau_{corr} < \tau_{gub} = \frac{1}{\omega_{min}} = 0.62 \left(\frac{100G}{B_o}\right)\left(\frac{n_o}{6x10^9 cm^{-3}}\right)^{1/2}\left(\frac{\ell}{10^9 cm}\right) s\ . \quad (67)$$

We have in this paper considered the possibility that second stage acceleration may occur during an/or immediately following the impulsive phase of proton flares. If primary energy release is associated with large amplitude Alfven waves self-consistently generated, then (66) admits the possibility of low harmonic standing waves playing a role in the basic flare instability. Phenomenologically, we may speculate that the correlation time (65) is related to the pulse width of hard x-ray bursts. Presently, observations show a fine structure to the bursts consisting of $\sim$ 1sec spikes [Svestka, 1976]. The Goddard x-ray spectrometer aboard the upcoming Solar Maximum Mission (SMM) satellite will have a temporal resolution $10^{-3}\ s < \tau < 10^{-1}\ s$ from 20-300 Kev energies. If, indeed, there is a relation between the Alfven wave correlation time and the hard x-ray pulse widths, we believe that the chances are good that this instrument will detect fine structure in the range of (65). The University of New Hampshire/Max Planck Institute gamma ray experiment also aboard SMM, with a temporal resolution down to 64 ms, should settle the question of how soon are > 10 Mev protons accelerated.

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FIGURE CAPTIONS

Figure 1. The temporal evolution of the hard x-ray and 2.2 Mev γ-ray bursts for the August 4, 1972 flare (from Lin and Hudson, 1976).

Figure 2. Resonant interaction of super-Alfvenic protons with a linear polarized elemental scattering center of length $x < x_{corr}$. The solid line portions of the trajectories indicate positions on this side of the page. Energization or deenergization depends on the cyclotron phase of the test particle. An average over a random-phased, isotropic distribution of particles leads to cyclotron damping of the oscillation.

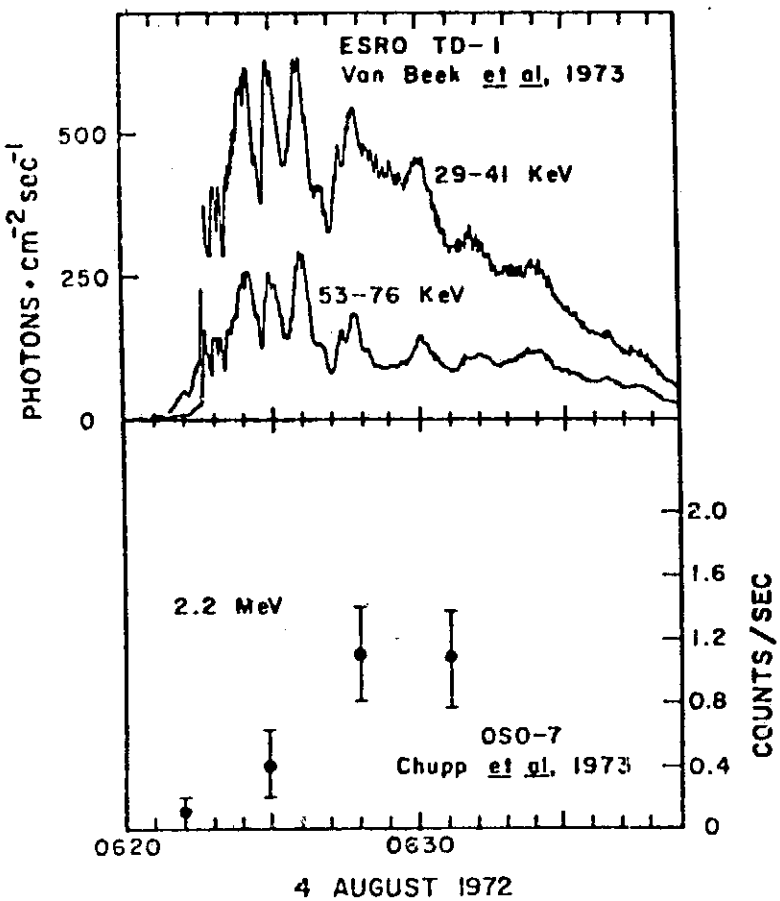

Figure 1

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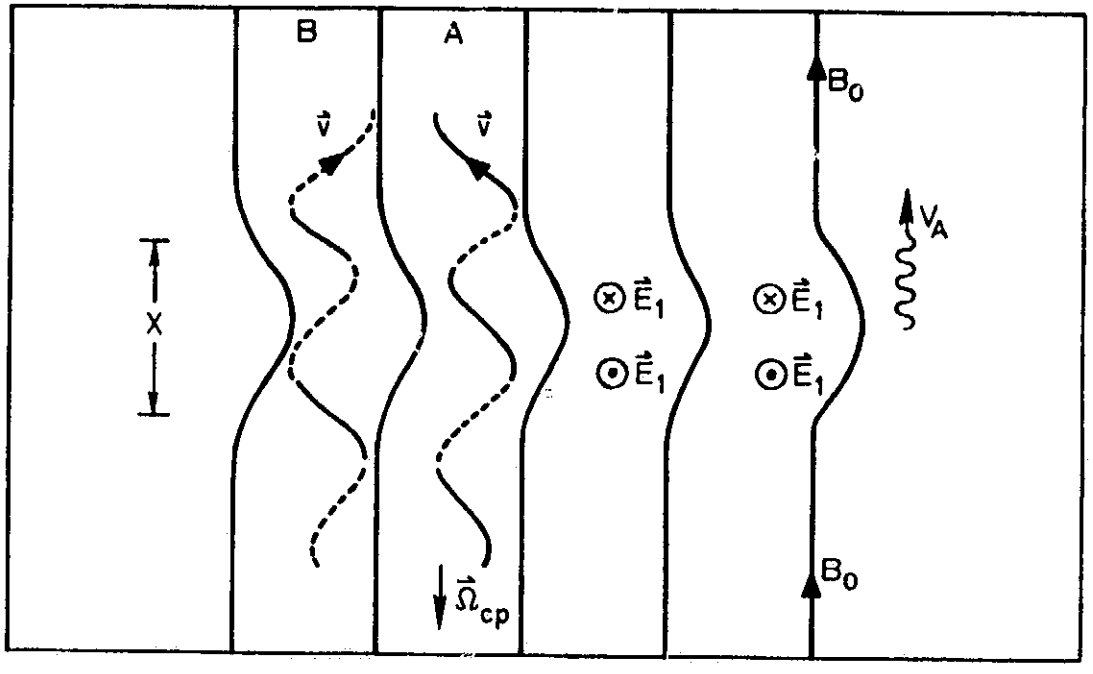
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Figure 2

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We consider the acceleration of solar flare protons by cyclotron damping of intense Alfven wave turbulence in a magnetic trap. The energy diffusion coefficient D_{EE} is computed for a near-isotropic distribution of super-Alfvenic protons and a steady-state solution for the particle spectrum is found for both transit-time and diffusive losses out of the ends of the trap. The acceleration time to a characteristic energy ∿20 Mev/nucl can be as short as 10sec. On the basis of phenomenological arguments we infer that the Alfven wave spectrum has a ω^{-2} frequency dependence and that the correlation time of the turbulence lies in the range 5×10^{-4}s $< \tau_{corr} < 5 \times 10^{-2}$ s.

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