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THE ACCELERATION OF ENERGETIC PARTICLES IN THE INTERPLANETARY MEDIUM BY TRANSIT-TIME DAMPING

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ABSTRACT

It has been reported recently by McDonald et al. (1976) that 1-MeV protons may undergo considerable acceleration in co-rotating streams. It has also been suggested recently by Fisk et al. (1974b) that interstellar neutral particles which are ionized in the solar cavity may be accelerated in the solar wind, and may account for the anomalous component that is observed in low-energy cosmic-rays (at ∿ 10 MeV/nucleon). It is shown here that the particles in both of these cases could be accelerated by transit-time damping propagating fluctuations in the magnitude of the interplanetary magnetic field (e.g. magnetosonic waves). The protons in co-rotating streams may be accelerated by transit-time damping the small-scale variations in the field magnitude that are observed at a low level in the inner solar system. The interstellar ions may be accelerated by transit-time damping large-scale field variations in the outer solar system.

The acceleration of energetic particles in the interplanetary medium has become a subject of renewed interest as the result of two recent observations. McDonald et al. (1976), by comparing Pioneer 11 and IMP 7 data, have found that the intensity of 1-MeV protons in co-rotating streams increases by more than an order of magnitude between 1 and 3 AU from the sun. Although other interpretations of this data may be possible, McDonald et al. (1976) draw the reasonable conclusion that the intensity increases are compelling evidence for extensive interplanetary acceleration of low energy particles. Further, the so-called anomalous component in low energy cosmic rays reported by McDonald et al. (1974) and Hovestadt et al. (1973) has been interpreted by Fisk, Kozlovoky, and Ramaty (1974b) as resulting from interstellar neutral particles which are swept into the solar cavity, and then ionized and accelerated in the solar wind. This interpretation, which is discussed in more detail in Fisk (1975 a, b, c), provides a natural explanation for the composition of the anomalous component, which is mainly He, N, O, and Ne. This interpretation requires, however, that the interstellar ions can be accelerated in the solar wind from energies $\sim$ 1 keV/nucleon, which is the energy they obtain immediately following their ionization, to their observed energy $\sim$ 10 MeV/nucleon.

A mechanism that is commonly invoked to accelerate particles in the interplanetary medium is a statistical Fermi process in which particles are resonantly scattered by Alfven waves with wavelengths comparable to the particle's gyro-radius (e.g. Jokipii, 1971). Equivalently, the particles are accelerated by cyclotron damping the Alfven waves. This mechanism, however, has two disadvantages. First, to provide a reasonable acceleration rate for 1 MeV protons, extensive pitch-angle scattering of the particles is required. For example, McDonald et al. (1976) find that even with unusually large Alfven speeds ($\sim$ 100 km/sec), the Fermi-scattering mechanism can account for the observed intensity increases only if the mean free path for 1-MeV protons near the orbit of earth and beyond is $\sim 7 \cdot 10^{-3}$ AU. In estimating this mean free path, McDonald et al. (1976) do not include the diffusion in energy that occurs in a statistical acceleration process; an effect which reduces the required acceleration rate. Also, as is noted by McDonald et al. (1976), the intensity increases appear to coincide with colliding streams in the solar wind, where conceivably the mean free path is unusually small. However, the required mean free path is more than an order of magnitude smaller that the typical value for low energy particles of $\sim$ 0.1 AU (Ma Sung et al., 1975). A second disadvantage of the Fermi-scattering mechansim is that its efficiency is reduced if the Alfven waves propagate primarily in one direction, as is frequently the case in the solar wind. However, in regions where streams interact, Alfven waves may be observed propagating in both directions along the magnetic field (Belcher and Davis, 1971).

As is discussed in more detail in Fisk (1975 c), arguements based on the expected modulation of galactic cosmic rays easily demonstrate that the Fermi-scattering mechanism is not responsible for the acceleration of the interstellar ions in the theory of Fisk et al. (1974b). In this theory the particles in the anomolous component are only singly-charged (during their dwell time in the solar cavity, the interstellar neutrals can be ionized only once (Fisk et al., 1974b)). Anomolous oxygen, then, in the energy range from 1 keV/nucleon to 10 MeV/nucleon (the range over which it must be accelerated) has a rigidity equal to that of, e.g., a proton in the energy range from 250 keV to 1.5 GeV, i.e. equal to the rigidity of modulated galactic protons. Since the mean free path is a function of rigidity, modulated cosmic rays and anomolous oxygen experience the same scattering conditions. If the scattering is sufficient to accelerate the anomolous oxygen by a Fermi-scattering process, then a modulation of galactic cosmic rays far in excess of allowable limits results.

In this paper, transit-time damping is examined as a possible means for accelerating low-energy particles in co-rotating streams, and interstellar ions in the theory of Fisk et al. (1974b). In this process particles are accelerated by interacting with propagating fluctuations in the magnitude of the interplanetary magnetic field (e.g. magnetosonic waves). This mechanism overcomes the two disadvantages associated with the Fermi-scattering process. As will be shown, this mechanism results in little pitch-angle scattering. Also, unlike Alfven waves, magnetosonic waves do not all have the same phase speed parallel to the mean magnetic field. Thus, there is an inherent dispersion to these waves, so that even if the waves propagate, e.g., primarily outward from the sun, a Fermi acceleration is still possible.

Magnitude fluctuations in the solar wind are present only at a relatively low level. The dominant fluctuations result from changes in the direction in the field that are caused primarily by Alfven waves. In the inner solar system, for example, the rms change in a component δB_x of the field, divided by the mean field strength squared, is $\langle (\delta B_x)^2 \rangle / B_o^2 \simeq 0.1 - 0.2$. In contrast, the rms change in the field magnitude is $\eta^2 \equiv \langle (\delta B)^2 \rangle / B_o^2 \simeq 0.01 - 0.03$ (Smith, 1974). The magnitude fluctuations are presumably reduced because these variations

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are damped not only by energetic particles, but also, as has been shown by Barnes (1966, 1968a,b, 1969) they are Landau damped by solar wind protons and transit-time damped by solar wind electrons. It is shown here, however, that the observed value of η^2 yields an acceleration rate for protons that is sufficient to account for the observations of McDonald et al. (1976). For the acceleration of interstellar ions in the outer solar system a larger value of η^2 will be required ($\eta^2 \approx 0.1$). However, this value will be justified on the grounds that the magnitude fluctuations involved are not readily damped by solar wind particles.

Magnitude fluctuations should be a continuous source of energy for accelerating energetic particles because these fluctuations are continuously regenerated. For example, Alfven waves decay into magnetosonic waves as they propagate through the spiral pattern of the interplanetary magnetic field (Belcher and Davis, 1971). Variations in the solar wind velocity result in compressions of the interplanetary field. Further, a stochastic magnetic field, like the interplanetary field, cannot be hydrostatically stable (Parker, 1972), and thus in continuously adjusting the particle and field pressure, should produce magnitude variations.

In section 2 the acceleration rates that are required to account for the observations of McDonald et al. (1976), and to account for the acceleration of interstellar ions in the theory of Fisk et al. (1974b), are determined. In section 3 the acceleration rates that result from particles transit-time damping magnitude fluctuations in the interplanetary magnetic field are calculated from standard quasi-linear/adiabatic theory, and are found to agree well with the required rates. The calculations presented in this paper provide formal justification for the acceleration rates used in Fisk (1975 c).

2. REQUIREMENTS FOR THE ACCELERATION

(2.1) Low-energy protons in co-rotating streams. Even with a mean free path as large as 0.1 AU, 1-MeV protons tend to co-move with the solar wind. An estimate of the acceleration rate that is required to account for the observations of McDonald et al. (1976) can then be obtained by considering that the particles move primarily by being convected with the solar wind, rather than by their own diffusive motion.

Consider that the acceleration is the result of a statistical Fermi process, or equivalently the result of the damping of a random distribution of waves. For example, the acceleration could be by the Fermi-scattering process (cyclotron damping), or it could be by transit-time damping. Such statistical acceleration can be described as a diffusion in momentum space, with a diffusion coefficient $\bar{D}_{pp} = \langle(\Delta p)^2\rangle / \Delta t$, which represents the rms change in the magnitude of particle momentum p, per unit time, averaged over all particle directions. The coefficient $\bar{D}_{pp}$ uniquely determines the rms change in energy and the mean change in energy per unit time; the latter term is the one considered by Jokipii (1971) and McDonald et al. (1976). In a steady-state model for a energetic particle stream in the interplanetary medium, the particle omnidirectional distribution function f (number of particles per unit volume of phase space, averaged over particle direction) behaves, then, according to the equation:

$$V \frac{\partial f}{\partial r} = \frac{1}{p^2} \frac{\partial}{\partial p} \left(p^2 \bar{D}_{pp} \frac{\partial f}{\partial p}\right) + \frac{p}{3r^2} \frac{\partial}{\partial r} (r^2 V) \frac{\partial f}{\partial p} \tag{1}$$

Here r is heliocentric radial distance and V is the solar wind speed. The first term on the right side of (1) describes the effects of the acceleration, and the second, the effects of adiabatic deceleration and convection in the solar wind (see Fisk et al., 1973). The differential

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intensity j, per unit interval of kinetic energy E, is related to f by $f = j/p^2$ (Forman, 1970).

Consider that $\bar{D}_{pp} = D_o p^{\beta}$, where D_o is a constant, independent of p and r. With the variable transformations $r \rightarrow r'$, $p \rightarrow p' r'^{-2/3}$, and then $r' \rightarrow ((7 - 2\beta)y/3)^{3/(7 - 2\beta)}$, (1) can be reduced to a simple diffusion equation of the type that is used in studying the propagation of energetic particles in solar flare events. Here y is the analogue of time in the flare problem, and p', the analogue of radial distance. For example, suppose that particles are injected into the accelerating region of the co-rotating stream at a position r_o and momentum p_o, such that r_o is much less than the position r at which observations are made, and p_o is much smaller than the observed momentum. Then, (1), written in terms of y and p', can be solved as is the flare propagation problem with impulsive injection of particles at the sun (Parker, 1963). The result in terms of r and p is

$$f(r,p) \propto r^{-(7-2\beta)} \exp\left[-\frac{1}{3}\frac{(7-2\beta)}{(2-\beta)^2}\frac{p^{2-\beta}V}{D_o r}\right] \qquad (2)$$

for $0 \leq \beta < 2$.

As is demonstrated below, β lies in the range 0-0.5 for transit-time damping. With $\beta \simeq 0$ and with a solar wind speed of V = 400 km/sec, (2) then predicts that a value of $\bar{D}_{pp}/p^2 = 1.5\cdot10^{-7}$ sec^{-1} for 1-MeV protons is required to account for the observations of McDonald et al. (1976). The intensity of these particles then increases by roughly a factor of ten between r = 1 and 3 AU. This value and form for $\bar{D}_{pp}$ yields, from (2), a spectral shape at r = 3 AU of $j \propto E \exp[-3.5 E]$, where E is in units of MeV. This spectral shape is not unrealistic, as can be seen by examining the proton intensities at different energies, reported in Figure 3 of McDonald et al. (1976).

(2.2) Interstellar ions. Interstellar neutral particles, unlike their ionized counterparts, are swept into the solar cavity by the motion of the sun through the interstellar medium. The neutrals, which should consist mainly of H, He, N, O and Ne, are then ionized here by photo-ionization from solar UV and by charge-exchange with the solar wind. (See Fisk et al. (1974b) for a defense of the neutral composition and Axford (1972) for a general review of this subject.) Once ionized, the interstellar particles, which are only singly-charged, are forced to co-move with the solar wind, and obtain energies in this process ~ 1 keV/nucleon (Fisk and Goldstein, 1976). In the theory of Fisk et al. (1974b) and Fisk (1975 a, c), it is assumed that a small fraction of the interstellar ions are then accelerated slowly but continuously in the solar wind, so that by the time the ions have been convected to large heliocentric distances (e.g. $r \sim 50$ AU), the particles have obtained energies ~ 10 MeV/nucleon. Singly-charged He, N, O, and Ne at these energies have high rigidities and can diffuse back into the inner solar system, producing the observed anomolous composition. Interstellar H, and any accelerated solar or galactic particles (which should be fully stripped of their electrons) are still at a low rigidity at these energies, and are excluded from the inner solar system by solar modulation (Goldstein et al., 1970; Gleeson and Urch, 1971).

A rough estimate for the spectral shape of the accelerated ions in the outer solar system can be obtained from (1) and (2), in which particles are assumed to move spatially by being convected by the solar wind, primarily, rather than by diffusing. This approximation is not unreasonable even at high rigidities. With relatively slow acceleration rates, the spatial gradients can adjust so that the loss of particles from the solar cavity is no larger by diffusion than it is by convection (Fisk, 1975 c). Detailed numerical calculations of the spectra of accelerated interstellar He and O, in which the effects of diffusion, convection, and adiabatic deceleration have been included, can be found in Fisk (1975 c).

Anomalous He particles are observed to obtain higher energies than do the anomolous N, O, or Ne (McDonald et al., 1974; Garcia-Munoz et al., 1973). Thus, the He observations place the most severe constraints on acceptable acceleration theories. Reproduced in Figure 1 is the He data for quiet-times in 1972, as was reported by Garcia-Munoz et al. (1973). Shown also in this figure is the spectrum that is expected for modulated galactic He. Clearly, to produce the flat portion of the spectrum in the energy range from $\sim$ 10-50 MeV/nucleon, which is the so-called anomolous portion, a spectrum of accelerated interstellar He at 1 AU, as is shown in Figure 1, is required.

With a value of $\bar{D}_{pp}/p^2 = 3 \cdot 10^{-8}$ sec^{-1} at E = 10 MeV/nucleon, and with β = 0.5 and V = 400 km/sec, (2) yields the spectral shape at r = 50 AU that is shown in the upper curve in Figure 1. The normalization of this curve is arbitrary since it depends critically on the exact number of interstellar ions that are actually accelerated. This number is expected to be only a small fraction ($\sim 10^{-4}$) of the number that are ionized in the solar cavity (Fisk, 1975 a, c). The modulation of the interstellar He, which is not excessive since the particles are singly-charged, should harden the spectrum as the particles propagate from r = 50 AU back into the inner solar system. Thus, this value of $\bar{D}_{pp}$ should be able to produce the required accelerated spectrum at 1 AU.

Note that the acceleration rate required for the interstellar ions is roughly equal to the rate required for protons in co-rotating streams. It was assumed that $\bar{D}_{pp}/p^2 \propto p^{-1.5}$. Thus, at 1 MeV/nucleon $\bar{D}_{pp}/p^2 = 1.7 \cdot 10^{-7} sec^{-1}$ for interstellar He, compared with $\bar{D}_{pp}/p^2 = 1.5 \cdot 10^{-7} sec^{-1}$ for the co-rotating protons. This comparison, however, is not necessarily meaningful since the acceleration in these two cases can occur under substantially different conditions. The acceleration in co-rotating streams is a local phenomenon in the inner solar system, which appears to occur near regions where solar wind streams interact. Interstellar ions are accelerated over much larger dimensions. To reach the required high energies of 10-50 MeV/nucleon, interstellar He must be accelerated continuously as it is convected out to distances of $r \sim 50$ AU.

3. ACCELERATION RATES

(3.1) <u>The formal calculation</u>. The acceleration rate $\bar{D}_{pp}/p^2$ that results from particles transit-time damping fluctuations in the magnitude of the interplanetary magnetic field can be calculated by applying the standard procedures of quasi-linear/adiabatic theory, as has been done in similar contexts by Barnes (1968a). Consider that the magnitude fluctuations result from the superposition of fast-mode magnetosonic waves. This mode is more likely to be present in the interplanetary medium than is, e.g., slow-mode waves (cf. Hollweg, 1975). In the hydro-magnetic limit the fluctuating magnetic field of a fast-mode wave must lie in a direction that is both normal to the wave vector $\underline{k}$ and normal to a unit vector $\hat{e}_1$ which lies in the $\underline{B}_o \times \underline{k}$ direction; $\underline{B}_o$ is the mean magnetic field (e.g. Lee and Völk, 1975). The fluctuating magnetic field at position $\underline{r}$ and time t is then

$$B'(\underline{r},t) = \int_{-\infty}^{\infty} d^3k \int_{-\infty}^{\infty} d\omega \exp\left[i\,(\underline{k}\cdot\underline{r} - \omega t)\right] \frac{A(\underline{k})}{k} (\underline{k} \times \hat{e}_1) \qquad (3)$$

Here, the function $A(\underline{k})$ is defined such that $A(\underline{k})A^*(\underline{k}) \equiv C(\underline{k})$, where the star denotes complex conjugate and

$$\int_{-\infty}^{\infty} d^3k \; C(\underline{k}) = \langle(\delta B)^2\rangle \tag{4}$$

with $\langle(\delta B)^2\rangle$ the rms change in the field magnitude, associated with the waves. From Maxwell's equations, the fluctuating electric field is

$$\underline{E}'(\underline{r},t) = \int_{-\infty}^{\infty} d^3k \int_{-\infty}^{\infty} d\omega \; \exp\,[i(\underline{k}\cdot\underline{r} - \omega t)] \; \frac{\omega A(\underline{k})\hat{e}_1}{ck} \tag{5}$$

where c is the speed of light. In this calculation only normal modes of the plasma are considered, which requires that the frequency ω satisfies the relationship

$$\omega^2 = k^2 \{1/2(V_A^2 + V_S^2) + 1/2[(V_A^2 + V_S^2)^2 - 4V_A^2 V_S^2 k_{\parallel}^2/k^2]^{1/2}\} \tag{6}$$

Here V_A and V_S are the Alfven and sound speeds, respectively, in the solar wind; $k_{\parallel}$ is the component of $\underline{k}$ parallel to B_o. The turbulence is also taken here to be homogeneous.

The fluctuating electric field will change the magnitude of the momentum of energetic particles. By following standard quasi-linear/adiabatic theory (cf. Goldstein et al., 1975), this change can be described by a diffusion coefficient.

$$D_{pp} = \frac{q^2}{c^2} \int_0^{\infty} dt \int_{-\infty}^{\infty} d^3k \; \exp\,[i(\underline{k}\cdot\underline{\tilde{r}} - \omega_o t)] \; \frac{C(\underline{k})\omega_o^2}{k^2} \; [\hat{e}_p\cdot\hat{e}_1]\,[\hat{e}_1\cdot\tilde{\hat{e}}_p] \tag{7}$$

Here q is particle charge and $\hat{e}_p$ is a unit vector in the direction of particle momentum $\underline{p}$. The tilde denotes streamed variables, which are to be evaluated along the trajectory of a particle spiralling along $\underline{B}_o$. Streamed variables are then functions of t.

After considerable manipulation, (7) can be reduced to

$$\frac{D_{pp}}{p^2} = \frac{1}{r_r{}^2 B_o{}^2} \int_0^{2\pi} d\psi \int_{-\infty}^{\infty} dk_{\parallel} \int_0^{\infty} k_{\perp} dk_{\perp}\ C\ (k_{\perp},\ k_{\parallel},\ \psi)\ \frac{\omega_o^2}{k^2}\ \mu'^2\ I \tag{8}$$

$$\text{where}\quad I = \pi \sum_{m=-\infty}^{\infty} \delta\ (k_{\parallel} v\mu - \omega_o + m\Omega)\ J_m{}'^2\ (\mu' k_{\perp} r_g) \tag{9}$$

Here $\underline{k}$ is defined in a cylindrical co-ordinant system, with the axis of symmetry $\underline{B}_o$; $k_{\parallel}$ and $k_{\perp}$ are measured parallel and normal to $\underline{B}_o$, respectively, and ψ is the azimuthal angle. Particle gyrofrequency in the mean field $\underline{B}_o$ is Ω, and particle Larmour radius is r_g. Particle speed is v and pitch angle μ, with $\mu' = (1 - \mu^2)^{1/2}$. The function J'_m (z) is the derivative with respect to z of a Bessel function of the first kind and order m; $\delta(z)$ is the delta function.

The diffusion coefficient for pitch angle scattering by magnetosonic turbulence, $D_{\mu\mu}$ $(=<(\Delta\mu)^2>/\Delta t)$, can be found by the same procedure to be

$$D_{\mu\mu} = \frac{q^2 \mu'^2}{c^2} \int_0^{2\pi} d\psi \int_{-\infty}^{\infty} dk_{\parallel} \int_0^{\infty} k_{\perp} dk_{\perp}\ C\ (k_{\perp},\ k_{\parallel},\ \psi)\ \frac{\left(\omega_o \mu/v - k_{\parallel}\right)^2}{k^2}\ I \tag{10}$$

The scattering rate $D_{\mu\mu}$ is roughly the inverse of the mean collision time, or equivalently it yields a mean free path $\lambda_{mfp} \sim v/D_{\mu\mu}$

A reasonable form for C $(k_{\perp}, k_{\parallel}, \psi)$ is

$$C\ (k_{\perp},\ k_{\parallel},\ \psi) = \frac{2\eta^2 B_o^2\ \Gamma(\nu + 2)\ \lambda_x \lambda_y \lambda_z}{3\pi^{3/2}\ \Gamma(\nu - 1/2)} \left[\frac{(\lambda_z^2 k_{\parallel}^2 + \lambda_{\perp}^2 k_{\perp}^2)}{(1 + \lambda_z^2 k_{\parallel}^2 + \lambda_{\perp}^2 k_{\perp}^2)^{\nu+2}}\right] \tag{11}$$

where $\lambda_{\perp}{}^2 = \lambda_x{}^2 \cos^2 \psi + \lambda_y{}^2 \sin^2 \psi$ and $\eta^2 = <(\delta B)^2>/B_o{}^2$. Here, λ_z, λ_y and λ_x are the correlation lengths for the turbulence in the direction parallel to $\underline{B}_o$ and the two orthogonal directions normal to $\underline{B}_o$, respectively; $\Gamma(z)$ is the gamma function. The index ν must exceed 0.5 so that the energy density in the field fluctuations is finite. If the field is sampled in, e.g., the direction along $\underline{B}_o$, this form for C yields a one dimensional power spectrum for fluctuations parallel to $\underline{B}_o$ which approaches a constant as $k_{\parallel} \to 0$ and declines as $k_{\parallel}^{-2\nu}$ for large $k_{\parallel}$. With ν in the range 0.5 - 1.5 this power spectrum resembles observed power spectra (Jokipii and Coleman, 1968; Sari and Ness, 1969, 1970).

The behavior of C near $\underline{k} = 0$ given in (11) is chosen so that the derivative of the correlation function in $\underline{k}$-space are well-defined at this point. Equivalently, the moments of the correlation function in $\underline{r}$-space are required to be finite. Dependences near $\underline{k} = 0$ in which C decreases less rapidly with $\underline{k}$, and for which the derivatives are not well defined, are also possible. In practice, the behavior of C near $\underline{k} = 0$ is relatively unimportant since it influences the acceleration rate of only quite high energy particles (i.e. $r_g >> \lambda_x, \lambda_y, \lambda_z$). The behavior of C near $\underline{k} = 0$ in (11) is in fact a conservative choice since it yields less acceleration for high energy particles than do other acceptable dependences.

It should be noted that in deriving (8) and (10) the particle distribution function was assumed to be gyrotropic about $\underline{B}_o$. This assumption is technically incorrect for turbulence that is not axial-symmetric about $\underline{B}_o$, i.e. it is incorrect when C varies with ψ. However, deviations from gyrotropy should occur only on times long compared with the time required for the integral over t in (7) to approximate closely its limiting value. In practice, the gyrotropic approximation should introduce no significant errors, even under the extreme circumstances considered below.

(3.2) <u>An illustration example.</u> Consider the case where $\nu > 1$ and $r_g << \lambda_x, \lambda_y$, and λ_z. This example is instructive in that it readily illustrates the physical processes involved in transit-time damping.

To a high accuracy only the m = 0 resonance need be considered in calculating D_{pp}/p^2 from (8). With $\nu > 1$ and $r_g << \lambda_x \lambda_y$, the Bessel function in this resonance can be well-approximated in (8) by $J_o'(\mu' k_\perp r_g) \approx -\mu' k_\perp r_g/2$. Further, to the accuracy required here it is reasonable to use the approximation that $\omega_o/k = \pm u$, where u is a constant phase speed. The weak dependence on ψ in the dispersion relation for fast-mode magnetosonic waves given in (6) is thus ignored. With this last approximation the resonance condition prescribed by the delta function in (8) requires (for m = 0) that

$$k_{\parallel} \mu v = \omega_o \simeq \pm uk \qquad (12)$$

which for energetic particles ($\mu v >> u$) yields

$$k_{\parallel} \simeq \pm \frac{u}{\mu v} k_\perp \qquad (13)$$

or thus $k_{\parallel} << k_\perp$. In (8), then, terms containing $k_{\parallel}$ can be neglected relative to terms in $k_\perp$, provided that the ratio of λ_z to λ_x and λ_y is such that $\lambda_z u/(\mu v \lambda_x)$ and $\lambda_z u/(\mu v \lambda_y)$ are also small.

With these approximations (11), (9), and (8) give

$$\frac{D_{pp}}{p^2} \simeq \frac{u^2 \mu'^4 \eta^2}{v|\mu|} \left[\frac{\Gamma(\nu+2)\, \lambda_x \lambda_y \lambda_z}{6\pi^{1/2}\, \Gamma(\nu - 1/2)}\right] \cdot \int_0^{2\pi} d\psi \lambda_\perp^2 \int_0^\infty dk_\perp k_\perp^5 (1 + \lambda_\perp^2 k_\perp^2)^{-\nu-2}, \quad \nu>1 \qquad (14)$$

which can be integrated to yield

$$\frac{D_{pp}}{p^2} \simeq \frac{u^2 \mu'^4 \eta^2}{v|\mu|} \left[\frac{\pi^{1/2}\, \Gamma(\nu-1)}{6\, \Gamma(\nu-1/2)}\right] \lambda_z \left[\frac{1}{\lambda_x^2} + \frac{1}{\lambda_y^2}\right] \qquad (15)$$

The terms in (15) can be readily identified, by the following argument.

A particle that interacts with a propagating fluctuation in the magnitude of the magnetic field experiences a change in the magnitude of its momentum, which can be expressed as

$$\frac{\Delta p}{p} \simeq \frac{u_{\parallel} \Delta v_{\parallel}}{v^2} \tag{16}$$

Here $u_{\parallel}$ is the phase speed of the fluctuation along the mean field, and $\Delta v_{\parallel}$ is the change in v parallel to the mean field, as viewed from the frame moving with $u_{\parallel}$. In deriving (16) use is made of the fact that the energy is constant in the frame moving with $u_{\parallel}$.

As can be seen from (14) and (13), with $\nu > 1$ the fluctuations that contribute most to D_{pp}/p^2 are those for which $k_{\parallel} = 1/\lambda_z' \ll 1/\lambda_z$ and $k_{\perp} \simeq (1/\lambda_x^2 + 1/\lambda_y^2)^{1/2}$. For such fluctuations $u_{\parallel}$ has a magnitude $u_{\parallel} \sim u\lambda_z'(1/\lambda_x^2 + 1/\lambda_y^2)^{1/2}$. Since there is no electric field in the frame moving with $u_{\parallel}$, $\Delta v_{\parallel}$ results from $\partial B/\partial z$, the gradient in the magnitude of the field, along B_o, which is produced by the fluctuation. With $r_g \ll \lambda_x$, λ_y and λ_z

$$\frac{dv_{\parallel}}{dt} \simeq -\frac{\mu'^2 v^2}{B_o} \frac{\partial B}{\partial Z} \tag{17}$$

or

$$|\Delta v_{\parallel}| \simeq \frac{\mu'^2 v^2 \eta \Delta t}{\lambda'_z} \tag{18}$$

Hence

$$\left|\frac{\Delta p}{p}\right| \simeq \mu'^2 u \eta \Delta t \left[\frac{1}{\lambda_x^2} + \frac{1}{\lambda_y^2}\right]^{1/2} \tag{19}$$

Here, $\Delta t \sim \lambda_z/(|\mu| v)$ is the time over which the particle interacts with a single, coherent fluctuation.

Clearly, to within a numerical constant D_{pp}/p^2 in (15) is $\sim$ $(\Delta p/p)^2/\Delta t$ where $\Delta p/p$ is given in (19).

In transit-time damping, then, particles are accelerated by interacting with a series of randomly moving gradients in the magnetic field. Although $\Delta v_{\parallel}$ is considered to be small in transit-time damping, this process is closely analogous with a classic Fermic process in which particles are accelerated by interacting with moving magnetic mirrors.

It should be noted that D_{pp}/p^2 in (15) varies inversely as the square of the correlation lengths normal to $\underline{B}_o$ and directly as the correlation length along B_o. Thus, the acceleration rate tends to be large when the characteristic length-scales of the turbulence are relatively small in general, or when the turbulence is elongated such that $\lambda_z >> \lambda_x$ and/or λ_y.

It should also be noted that particles interact here with long-wavelength waves. Thus, the magnetic moment is a constant of the motion, as was assumed in deriving (17) and (18). Transit-time damping, then, can increase the energy of a particle only in motion parallel to $\underline{B}_o$, or equivalently transit-time damping decreases the average particle pitch-angle. The acceleration rate, however, varies as μ'^4. For the acceleration to remain efficient, then, pitch-angle scattering must isotropize the particle distribution on a time-scale of at least p^2/D_{pp}. The pitch-angle scattering that normally occurs in the interplanetary medium ($\lambda_{mfp} \sim 0.1AU$) is generally more than sufficient for this purpose.

(3.3) A more realistic example. One-dimensional power spectra that are observed in interplanetary space generally fall off less steeply than k^{-2} for large k, which implies that $\nu<1$ (Jokipii and Coleman, 1968; Sari and Ness, 1969, 1970). With such values for ν, the integral over

$k_\perp$ in (14) does not exist. The approximation $J_o'(\mu' k_\perp r_g) \simeq -\mu' k_\perp r_g/2$, which was used in deriving (14), is not valid at large values of $k_\perp$, and with $\nu<1$ causes the integral to diverge.

An accurate value for D_{pp}/p^2 when $\nu<1$ can be obtained by integrating (8) and (9) numerically, as is done below. An analytic approximation to D_{pp}/p^2 can also be obtained by noting that the Bessel functions in (9), rather than causing (8) to diverge, actually truncate the integral over $k_\perp$ at values of $k_\perp > 1/\mu' r_g$. This truncation is to be expected because the effects of fluctuations that have scale-sizes small compared with $\mu' r_g$ should average to zero, over a gyro-period. To a crude approximation, then, D_{pp}/p^2 can be estimated by letting $J_o'(\mu' k_\perp r_g) \simeq -\mu' k_\perp r_g/2$ and integrating only over the range $0 \le k_\perp \le 1/\mu' r_g$, or

$$\frac{D_{pp}}{p^2} \simeq \frac{u^2 \mu'^4 \eta^2}{v|\mu|} \left[\frac{\Gamma(\nu+2)\lambda_x \lambda_y \lambda_z}{6\pi^{1/2}\Gamma(\nu-1/2)}\right] \int_0^{2\pi} d\psi \lambda_\perp^2 \int_0^{1/\mu' r_g} dk_\perp k_\perp^5 (1+\lambda_\perp^2 k_\perp^2)^{-\nu-2}, \quad \nu<1 \qquad (20)$$

With $\mu' r_g << \lambda_\perp$, (20) yields a substantially larger acceleration rate than does (14). This increased acceleration can be attributed to the smaller length-scales associated with the turbulence when $\nu<1$. As can be seen in (19), for example, the momentum change of a particle interacting with a wave varies directly as the amplitude of the wave and inversely as the wavelength normal to $\underline{B}_o$. With $\nu>1$, this ratio is simply δB over the perpendicular correlation length $(1/\lambda_x^2 + 1/\lambda_y^2)^{-1/2}$. With $\nu<1$, however, there is sufficient energy in the smaller wavelength fluctuations to yield a ratio large compared with $\delta B(1/\lambda_x^2 + 1/\lambda_y^2)^{1/2}$.

It should be noted that D_{pp}/p^2 in (20) is now a function of r_g. For $r_g << \lambda_x$ and λ_y, it can be shown that D_{pp}/p^2 varies as $r_g^{-2(1-\nu)}$

($\nu<1$). The momentum dependence is $D_{pp}/p^2 \propto p^{-3+2\nu}$.

When $\mu' r_g >> \lambda_x$ and λ_y, D_{pp}/p^2 in (20) can be integrated to yield

$$\frac{D_{pp}}{p^2} \simeq \frac{u^2 \mu'^4 \eta^2}{v|\mu|} \left[\frac{\pi^{1/2}\Gamma(\nu+2)}{36\Gamma(\nu-1/2)}\right] \lambda_z \left[\frac{1}{\lambda_x^2} + \frac{1}{\lambda_y^2}\right] \left(\frac{\lambda_x \lambda_y}{\mu'^2 r_g^2}\right)^3 \quad (21)$$

Upon comparing (21) with (15), it can be seen that the acceleration is inefficient when $\mu' r_g >> \lambda_x$ and λ_y.

(3.4) Comparison with required acceleration rates. To compare the acceleration rate for transit-time damping, with the required rates discussed in Sections 2.1 and 2.2, (8) and (9) must be averaged over particle pitch-angle to form $\bar{D}_{pp}/p^2$. A straight forward average

$$\bar{D}_{pp} = \int_{-1}^{1} d\mu \, D_{pp} \quad (22)$$

however, does not seem to be a meaningful procedure. As can be seen in (15), (20) and (21) D_{pp}/p^2 becomes large as $\mu \to 0$, and thus $\bar{D}_{pp}$ in (22) depends strongly on the behavior of D_{pp} at small values of μ. Unfortunately, at small μ quasi-linear/adiabatic theory is unlikely to be correct (cf. Fisk et al., 1974a).

It should be noted, however, that the integral in (22) does exist, despite the apparent divergence of D_{pp} as $\mu \to 0$. The approximations $\lambda_z u/(\mu v \lambda_x) << 1$ and $\lambda_z u/(\mu v \lambda_y) << 1$, which were used in deriving (15), (20) and (21) fail for small μ, and in fact if computed accurately, $D_{pp} = 0$ at $\mu = 0$.

Rather than performing an average, the expedient and conservative procedure that is followed here is simply to take $\bar{D}_{pp} = D_{pp}$, where D_{pp} is evaluated at a characteristic value of $\mu = 0.3$.

(3.4.1) Low-energy protons in co-rotating streams. The co-rotating proton streams reported by McDonald et al. (1976) appear to coincide with stream-

stream interaction regions in the solar wind. In such regions it is conservative to take $\eta^2 = 0.01$, $u = 75$km/sec and $\lambda_x = \lambda_z = 2 \cdot 10^{10}$cm. This value for η^2 is typical of interplanetary conditions in the inner solar system (Smith, 1974). Similarly, in specifying u it is assumed simply that the Alfven speed and ion thermal speed are equal at about 50 km/sec, which is again characteristic of average conditions. The correlation lengths used here are equal to the value found by Fisk and Sari (1973) for small-scale propagating waves in the solar wind. This correlation length is smaller than the value of $2 \cdot 10^{11}$ cm given by Jokipii and Coleman (1968); however, Fisk and Sari (1973) argue that the length $2 \cdot 10^{11}$cm applies only to large-scale tangential discontinuities in the interplanetary field. Further, it seems reasonable to use a small correlation length for an interaction region, in which the turbulence may be generated locally.

Upon assuming that the particles are protons, that $B_o = 5 \cdot 10^{-5}$Gauss, and that $\nu = 0.75$ and $\mu = 0.3$, (8) and (9) have been integrated numerically with the above values for η^2, u, and λ_x, λ_y and λ_z. Several resonances were computed, but again only the m = 0 resonance is important. Shown in Figure 2 is $\bar{D}_{pp}/p^2$ as a function of energy. It can be seen here that $\bar{D}_{pp}/p^2 \simeq 1.2 \cdot 10^{-7}$ at 1 MeV, which is close to the acceleration rate that is required to account for the intensity increases reported by McDonald et al. (1976) (cf. Section 2.1).

It should be noted in Figure 2 that $\bar{D}_{pp}/p^2$ varies roughly as p^{-2}. This dependence is intermediate between $\bar{D}_{pp}/p^2 \propto p^{-1.5}$, which holds when $r_g \ll \lambda_x, \lambda_y$, and $\bar{D}_{pp}/p^2 \propto p^{-7}$ which, from (21), is the case when $r_g \gg \lambda_x, \lambda_y$. The dependence $\bar{D}_{pp}/p^2 \propto p^{-2}$ yields $\beta \simeq 0$ ($D_{pp} = D_o p^{\beta}$), which is the value used in the discussion in Section 2.1.

Also plotted in Figure 2 is the scattering mean free path $\lambda_{mfp} \sim v/D_{\mu\mu}$, where $D_{\mu\mu}$ is calculated from (10) by using the same parameters as were used in computing $\bar{D}_{pp}/p^2$. The most important contribution in this case comes from the m = 1 resonance, i.e. the scattering is due mainly to fluctuations with scale-sizes comparable to the particle gyroradii. The mean free path here is larger than the typical value of 0.1 AU, and thus acceleration by transit-time damping does not require or imply extensive pitch-angle scattering. Presumably, pitch-angle scattering in the interplanetary medium is the result primarily of particles interacting with Alfven waves, which are not considered here.

(3.4.2) Interstellar ions. As was discussed in Section 2.2, the acceleration of interstellar He ions in the solar wind is assumed to occur continuously so that as the particles are convected to large heliocentric distances they obtain energies of $\gtrsim$ 10 MeV/nucleon. In the outer solar system, the magnetic field is relatively weak and thus the gyro-radius of a 10-MeV/nucleon particle is large. For example, if the field strength falls off as 1/r, which is the case if the field pattern is that of an Archimedes spiral, then at r = 50 AU, $B_o \sim 10^{-6}$ Gauss, and r_g for a singly charged, interstellar He ion at 10-MeV/nucleon is $\sim$0.1AU. As can be seen from (20), however, acceleration by transit-time damping is inefficient when $r_g \gg \lambda_x, \lambda_y$. Thus, acceleration in the outer solar system can result only from the interaction of particles with large-scale field structures.

In addition to small-scale turbulence, such as is used above in accelerating co-rotating protons, stream-stream interactions in the solar wind may also produce the required large-scale variations in the interplanetary magnetic field. As a fast stream overtakes a slow one, the solar wind density, pressure, and magnetic field tend to be compressed throughout the interaction region. As is discussed in detail by, e.g., Gosling et al. (1975), within the first few AU from the Sun these compressions become enhanced with increasing heliocentric distance, and eventually steepen into shocks. The compressions at this point occupy isolated regions of space, a result which is compatible with the Pioneer 10 observations (Smith and Wolf, 1975). The shocks, however, propagate relative to the solar wind and eventually should cause the compressions to spread and thus to disipate. This spreading could conceivably cause adjacent compression regions to merge with each other. Thus, at large heliocentric distances, e.g. $r>10$ AU, the compression regions could conceivably disipate into a homogeneous distribution of propagating large-scale fluctuations in the magntiude of the interplanetary field.

These large-scale fluctuations should have a highly elongated shape in the outer solar system; a configuration, which as was discussed in Section 3.2, should make them efficient accelerators. The solar wind expands in the azimuthal and polar directions about the Sun, but does not expand appreciably in the radial direction. With the field in an Archimedes spiral pattern, λ_z is measured in essentially the azimuthal direction at large r; λ_x is taken to lie in the polar direction, and then λ_y in the radial direction. Thus, λ_z and λ_x both increase in proportion to r, with λ_y held constant.

Consider that $\eta^2 = 0.1$, $u = 75$km/sec, $\nu = 0.75$ and $\mu = 0.3$. The correlation lengths are taken to have the above radial dependences, and at $r = 1$ AU to have the values $\lambda_y = 0.2$ AU and $\lambda_x = \lambda_z = 0.4$ AU. The mean field is assumed to be $B_o = 5\cdot10^{-5}$ Gauss at $r = 1$ AU, and to decrease as $1/r$. With these values and upon assuming that the particles are singly charged He ions, $\bar{D}_{pp}/p^2$ has been calculated by integrating (8) and (9) numerically.

Plotted in Figure 3 is $\bar{D}_{pp}/p^2$ as a function of energy at various heliocentric distances. It can be seen here that at r = 50 AU and at an energy of 10 MeV/nucleon $\bar{D}_{pp}/p^2 \approx 3 \cdot 10^{-8} sec^{-1}$. As was discussed in section 2.2, this rate is sufficient to account for the acceleration of interstellar ions in the theory of Fisk et al. (1974b) and Fisk (1975c). At energies above $\sim$ 1 MeV/nucleon, $\bar{D}_{pp}/p^2$ varies roughly as $p^{-1.5}$. Thus, $\bar{D}_{pp} = D_o p^{0.5}$, which is the functional form used in Section 2.2.

The value of $\eta^2 = 0.1$ used here is larger than the value of η^2 = 0.01 - 0.03 which is observed in the inner solar system (Smith, 1974). Further, η^2 is taken to be constant here, independent of r, and thus no consideration has been given to the possibility that η^2 may decrease with r as the waves do work on the solar wind (Hollweg, 1974). The constant value of $\eta^2 = 0.1$, however, could be justified in the outer solar system, since large-scale elongated fluctuations may not be readily damped by solar wind particles. As can be seen in Figure 3, at large heliocentric distances the acceleration rate for 1-kev/nucleon ions becomes small. In this region, the fluctuations have become sufficiently elongated so that for these low-energy ions the ratio $\lambda_z u/(\lambda_y |\mu| v) >> 1$. Thus, the transit-time of a particle over a correlated fluctuation $(\lambda_z /|\mu| v)$ is long compared with the characteristic time for the fluctuation to change λ_y/u. The particle gains only a small portion of the energy that the gradient in the field strength can impart, before the gradient changes, and thus the acceleration is reduced. The speed of, e.g., solar wind electrons is comparable to the speed of 1-kev/nucleon ions, and thus the acceleration of, or the damping by, solar wind particles should be similarly reduced.

The interplanetary magnetic field should lie in the azimuthal direction in the outer solar system only in some average sense. Significant local deviations from this direction are expected. In calculations of the behavior of the scattering mean free path with heliocentric distance, these local deviations have an appreciable effect. For example, Völk et al. (1974) found that if the field lies strictly in the azimuthal direction, the spatial diffusion coefficient increases rapidly beyond earth and becomes imcompatible with modulation observations. In a subsequent paper, however, Morfill et al. (1975) found by including local deviations that the diffusion coefficient is consistent with, for example, measurements of the cosmic-ray gradient from Pioneer 10.

In contrast, local deviations in the field direction do not appear to have an appreciable effect on the acceleration rate from transit-time damping. As can be seen from (19) the momentum change of a particle interacting with a propagating fluctuation in the field magnitude depends on the magnitude of the $\underline{k}$-vector of the fluctuation, not on its direction relative to $\underline{B}_o$. In the case considered in (19), $|k| \sim (1/\lambda_x^2 + 1/\lambda_y^2)^{1/2}$. Thus, the field can undergo small-scale changes in direction within a correlated segment of turbulence (these directional changes could be the result of, e.g., Alfven waves) without appreciably affecting the total momentum change. The insensitivity of D_{pp}/p^2 to the local field direction can also be seen in (14) and (20). The integral in these expressions depend only weakly on $k_{\parallel}$.

CONCLUDING REMARKS

It was shown in this paper that acceleration by transit-time damping can account for the intensity increases in co-rotating proton streams that have been reported by McDonald et al. (1976). The small-scale fluctuations in the magnitude of the interplanetary magnetic field that are observed in average conditions in the inner solar system appear to yield a sufficient acceleration rate. It should be noted here, however, that in these calculations the waves were assumed to be fast-mode magnetosonic waves and the correlation length was taken to be reasonable, but small. At times when the fluctuations are caused by other modes or the correlation length is larger, the acceleration rate could be less. Indeed the coincidence of the events observed by McDonald et al. (1976) with stream-stream interaction regions in the solar wind may result because fast-mode magnetosonic waves and small correlation lengths are characteristic of such interaction regions.

It was also shown here that interstellar particles which are ionized in the solar cavity can be accelerated to energies $\gtrsim$ 10 MeV/nucleon by transit-time damping large-scale variations in the magnitude of the interplanetary mangetic field in the outer solar system. These interstellar particles can thus account for the anomolous component observed in low-energy cosmic-rays, as has been proposed by Fisk et al. (1974b). It is suggested here that the required large-scale field variations could be the remnants of stream-stream interaction region in the solar wind. Of course, this is only speculation, and only with direct observation or with more detailed models of the evolution of solar wind streams will it be possible to determine whether the required fluctuations exist.

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FIGURE CAPTIONS

Figure 1. The differential intensity of accelerated interstellar He that is required at 1 AU. This intensity when combined with the expected flux of galactic He produces the anomolous or flat portion of the observed spectrum. Also shown is the intensity of accelerated interstellar He at 50 AU, which is predicted from (2).

Figure 2. A plot vs. kinetic energy of the acceleration rate $\bar{D}_{pp}/p^2$ that results from protons transit-time damping fluctuations in the magnitude of the interplanetary magnetic field. The parameters used in calculating this curve should describe the small-scale fluctuations in the field magnitude which occur in the inner solar system, in stream-stream interaction regions of the solar wind. Also shown is the characteristic mean free path for pitch-angle scattering which results from the magnitude fluctuations that are used in calculating $\bar{D}_{pp}/p^2$.

Figure 3. A plot vs. kinetic energy, at various radial distances, of the acceleration rate $\bar{D}_{pp}/p^2$ that results from singly-charged interstellar helium ions transit-time damping fluctuations in the magnitude of the interplanetary magnetic field. The parameters used in calculating these curves may describe large-scale variations in the field magnitude in the outer solar system.

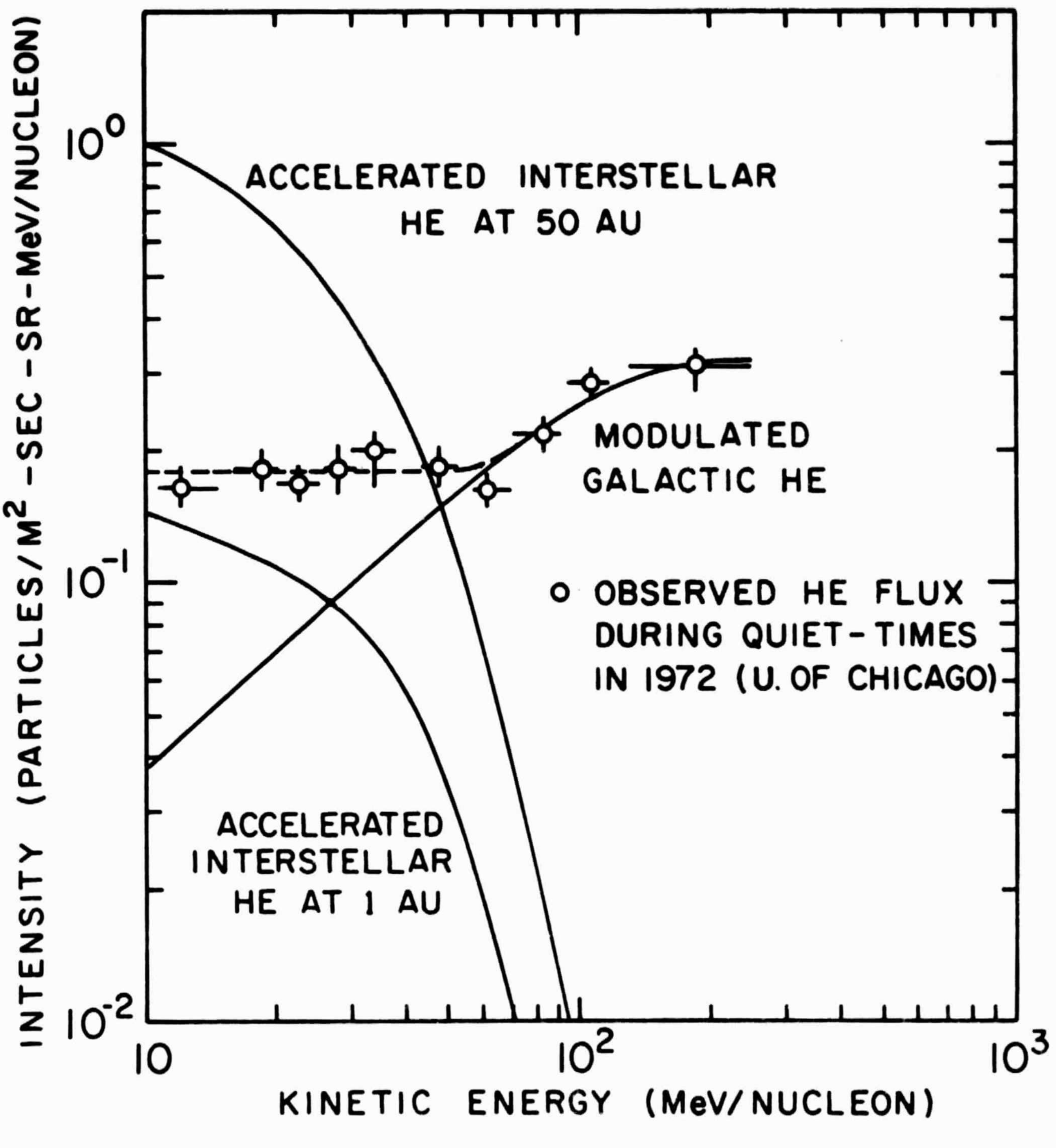

Figure 1

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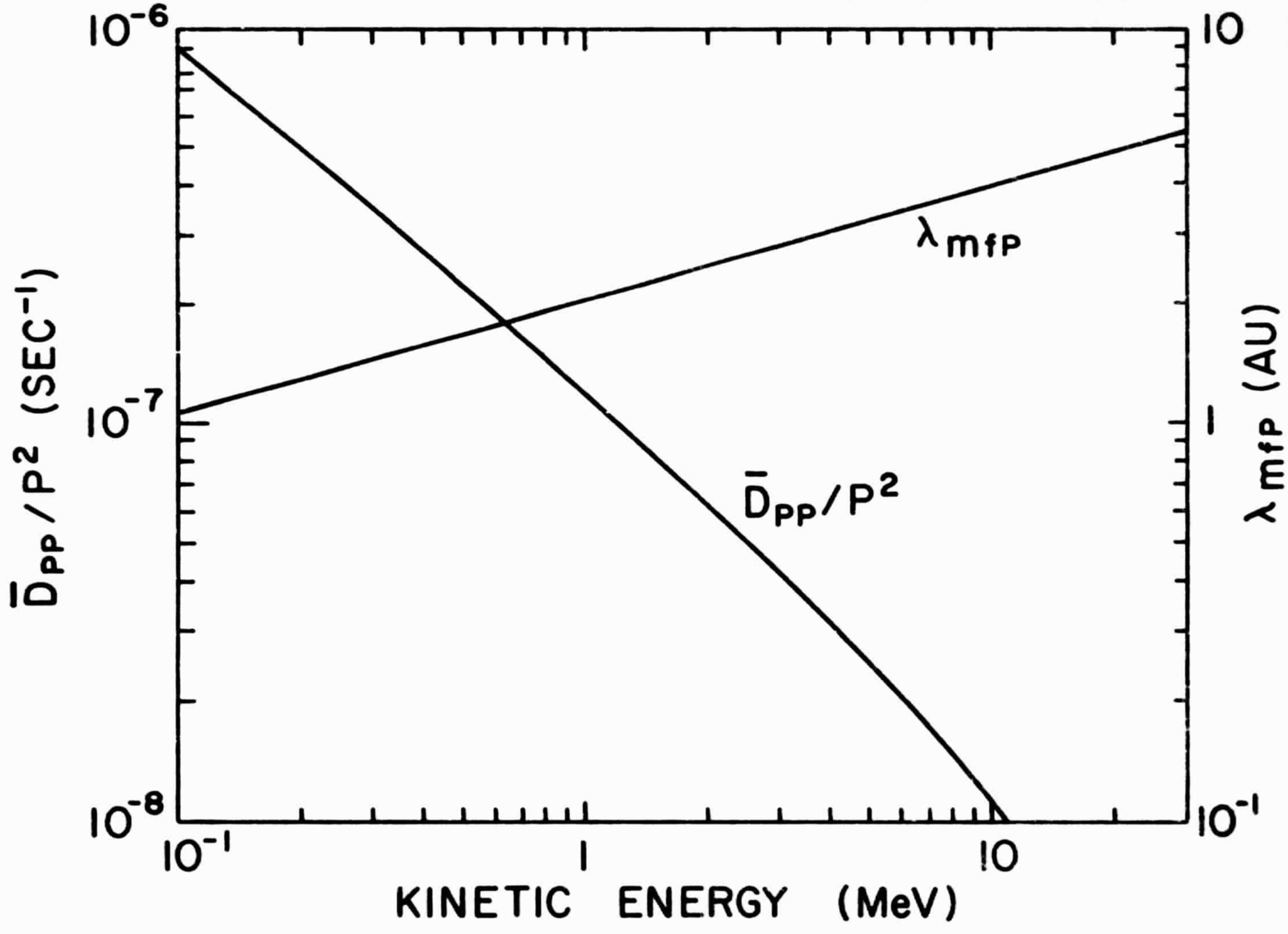
10^-6
10^-7
10^-8
10^-1
1
10
10
1
10^-1
λ mfP
D̄PP/P²
D̄PP/P² (SEC^-1)
λ mfP (AU)
KINETIC ENERGY (MeV)

Figure 2

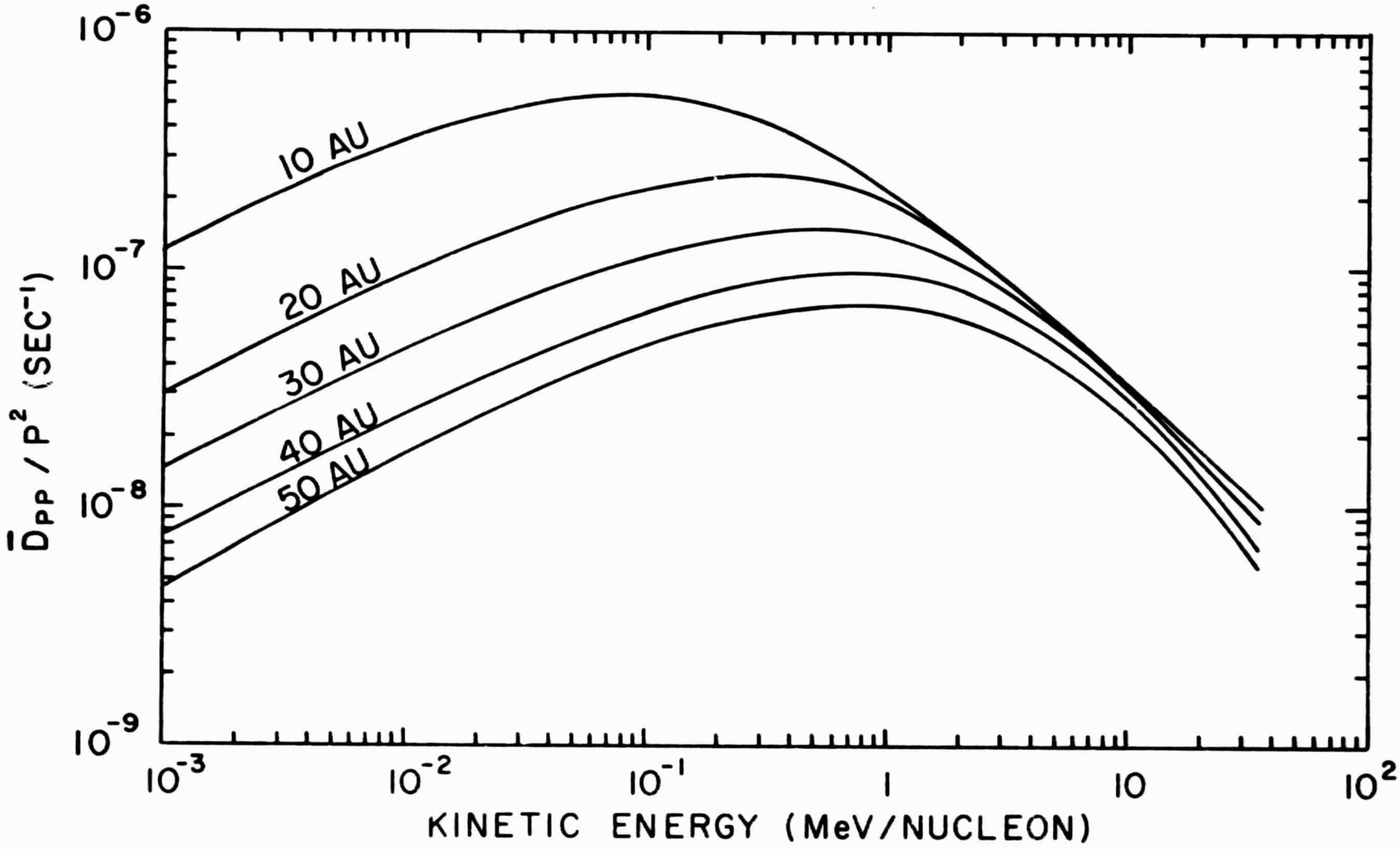

Figure 3

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