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RELATIVISTIC ELECTRONS

AND

WHISTLERS IN JUPITER'S MAGNETOSPHERE

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ABSTRACT

We have computed the path-integrated gain of parallel propagating whistlers driven unstable by an anisotropic distribution of relativistic electrons in the stable trapping region of Jupiter's inner magnetosphere. The requirement that a gain of 3 e-foldings of power balance the power lost by imperfect reflection along the flux tube sets a stably-trapped flux of electrons $J^* = 4 \times 10^{10} L^{-4} cm^{-2} sec^{-1}$ which is close to the non-relativistic result. Comparison with measurements shows that observed fluxes are near the stably-trapped limit, which suggests that whistler wave intensities may be high enough to cause significant diffusion of electrons accounting for the observed reduction of phase space densities. A crude estimate of the wave intensity necessary to diffuse electrons on a radial diffusion time scale yields a magnetic field fluctuation intensity of $I_B = 1.5 \quad 10^{-18} (\Omega_{ce}/\omega)^2 L^{\alpha-2}$ watts $m^{-2} Hz^{-1}$ as a lower limit.

I INTRODUCTION

Decreases in observed phase space densities of energetic electrons in Jupiter's inner magnetosphere have led to the suggestion that pitch angle scattering of the particles by wave microturbulence is an active loss mechanism (Fillius et al., 1975; McIlwain and Fillius, 1975; Baker and Van Allen, 1976). A likely suspect is the electron whistler wave or R mode, the theory of which when applied to earth's environment has been successful in explaining electron precipitation losses in the stable trapping region of the inner magnetosphere (Kennel and Petschek, 1966; Lyons et al., 1972).

Radio observations at decimeter wavelengths of the synchrotron emissions from Jupiter's radiation belts indicated long before Pioneer 10 that the electron distribution was highly anisotropic and that most of the energetic electrons were confined to the magnetic equator by a pancaked pitch angle distribution (Roberts and Komesaroff, 1965; Thorne, 1965). Pioneers 10 and 11 confirmed this expectation and also verified the theory that the immediate source of the radiation belt electrons was inward radial diffusion which, conserving the first and second particle adiabatic invariants, would flatten the pitch angle distribution. Since induced emission of whistlers is a consequence of anisotropic distribution, the hypothesis that whistler turbulence is responsible for observed losses is an attractive one both because of the ample growth rates possible and the theoretical simplicity of the instability.

The first hint that non-synchrotron associated losses were present at Jupiter was given by Stansberry and White (1974) before Pioneer 10. They set up a radial diffusion model for electrons to compute the strip-scan brightness and flux density spectrum of synchrotron emissions and found that an *ad hoc* loss process was needed at low L values for a reasonable fit with radio observations. Coroniti (1974), in deriving a comprehensive theoretical model of radiation belt electron fluxes, employed the stably-trapped limit concept for whistlers using the best pre-Pioneer 10 values for relevant parameters. The data from Pioneers 10 and 11 have narrowed the range of parameter space (in particular, the cold plasma density and anisotropy) available for theoretical models and the observed losses invite a reexamination of the relativistic electron-whistler interaction.

The purpose of this paper is to detail some of the consequences of relativistic electrons in stably-trapped equilibrium with parallel-propagating whistler waves. Approximate scaling laws for the stably-trapped electron flux and equilibrium wave intensity are derived. For simplicity of analysis and clarity of content, the major restrictions to our model are the following: 1) we treat the waves as generated locally and travelling strictly parallel to the ambient dipole magnetic field; 2) all resonant electrons are ultrarelativistic in that the total energy is proportional to particle momentum, $E \simeq pc$.

In Section II the equatorial growth rate for whistlers is derived for a distribution modeled as $f(\vec{p}) \propto p^{-(N+2)} \sin^M \theta$.

The logarithmic gain for maximally-amplified waves is computed by an approximation to the path-integrated growth rate along a flux tube and the stably-trapped limit is defined as that which will produce 3 e-foldings of power. A similar model for the distribution was used by Liemohn (1967) to compute whistler amplification by relativistic electrons at earth. Schulz and Vampola (1975) considered an artificial radiation belt produced by the beta decay of nuclear-fission debris and computed the relativistic stably-trapped limit for a distribution with an angular distribution similar to ours but with an exponential energy dependence. Our procedure and results are consistent with the above authors. The stably-trapped limit of relativistic electron fluxes is found to compare closely with non-relativistic Kennel-Petchek theory. Comparison with measurements shows that observed fluxes lie near the limit, which lends support to the idea that whistlers are active. In Section III we briefly treat the aspects of quasilinear diffusion of relativistic electrons and estimate a level of wave intensity that will support diffusion (losses) on a time scale comparable to radial diffusion.

II RELATIVISTIC ELECTRON-WHISTLER INTERACTION

a. Equatorial Growth Rate

If we restrict our attention to whistlers propagating parallel to the local magnetic field $\vec{B}_o = B_o\hat{z}$, the dispersion relation for R modes with $k_\perp = 0$ is (Lerche, 1969)

$$K(k,\ \omega) = 0 = -\frac{k^2c^2}{\omega^2} + 1 + \sum \frac{2\pi e^2}{m\omega}\int d\vec{p}\,\frac{p_\perp}{\omega\gamma - kp_z/m + \Omega} \cdot \left\{\frac{\partial f_o}{\partial p_\perp} + \frac{k}{m\omega\gamma}\left[p_\perp \frac{\partial f_o}{\partial p_z} - p_z \frac{\partial f_o}{\partial p_\perp}\right]\right\} \qquad (1)$$

The $\sum$ represents a sum over particle species (species labels on all quantities are implicit), $\gamma = \left[1 + p^2/m^2c^2\right]^{\frac{1}{2}}$, m is the particle rest mass, and $\Omega = \frac{qB_o}{mc}$. If the particle distribution function $f_o(\vec{p})$ is composed of a cold plasma background and a small relativistic component f, then defining $t = p/mc$ and $x = \cos\theta = \frac{\vec{p}\cdot B_o}{pB_o}$, equation (1) can be written in spherical coordinates, assuming azimuthal symmetry in $\vec{p}$-space, as

$$\frac{k^2c^2}{\omega^2} = 1 + \frac{\omega_p^2}{\left(\Omega_{ci} + \omega\right)\left(\Omega_{ce} - \omega\right)} - \sum \frac{4\pi^2e^2}{m\omega kc} \cdot \int_0^\infty t^2dt \int_{-1}^{1} \frac{\left(1 - x^2\right)dx}{x - \left(\frac{\omega\gamma + \Omega}{kct}\right)}\left\{\frac{\partial f}{\partial t} + \frac{\partial f}{\partial x}\left(\frac{kc}{\omega\gamma} - \frac{x}{t}\right)\right\} \qquad (2)$$

where

$$\omega_p^2 = \frac{4\pi n_o e^2}{m} \quad \text{and} \quad \Omega_{ce} = \left|\frac{eB_o}{mc}\right|$$

for the electrons. If the wave number has a small imaginary part, $k = k_r + ik_i$, we can approximate $k_i \approx \frac{-K_i}{\partial K_r/\partial k}$ and integrating around the indented pole of the integrand for the electron contribution yields

$$k_i(\omega) = \frac{-2\pi^3 e^2 \omega}{mk^2c^3}\int_{t_o}^{\infty} dt\ t^2\left(1 - x^2\right)\left[\frac{\partial f}{\partial t} + \frac{\partial f}{\partial x}\left(\frac{n}{\gamma} - \frac{x}{t}\right)\right]\Bigg|_{x = \frac{\omega\gamma - \Omega_{ce}}{kct}} \tag{3}$$

The integration is taken along the resonance contour $x = \frac{\omega\gamma - \Omega_{ce}}{kct}$ and t_o represents the minimum momentum (magnitude) satisfying $-1 = \frac{\omega\gamma_o - \Omega_{ce}}{kct_o}$. The subscript on the real part of the wavenumber has been dropped $\left[k_r \rightarrow k\right]$ and

$$\frac{k^2c^2}{\omega^2} = n^2 = 1 + \frac{\omega_p^2}{\left(\Omega_{ci} + \omega\right)\left(\Omega_{ce} - \omega\right)} \tag{4}$$

The resonance curve in momentum space is a hyperbola, plotted in Fig. 1, satisfying the relation

$$\left[t_z + \frac{n\bar{\Omega}}{n^2 - 1}\right]^2 - \frac{t_\perp^2}{n^2 - 1} = \frac{\bar{\Omega}^2}{\left(n^2 - 1\right)^2} + \frac{1}{n^2 - 1} \tag{5}$$

with $\bar{\Omega} = \frac{1}{\bar{\omega}} = \frac{\Omega_{ce}}{\omega}$. The minimum momentum is given by

$$t_o = \frac{n\bar{\Omega} - \sqrt{\bar{\Omega}^2 + n^2 - 1}}{n^2 - 1} \tag{6}$$

The resonance curve crosses the $t_\perp$ axis at the point $t_\perp = \sqrt{\bar{\Omega}^2 - 1}$ and asymptotes to the angle $\varphi = \sin^{-1}\frac{1}{n}$. The passage to the non-relativistic regime requires $n \gg 1$, $t \ll 1$ and the resonance curve flattens to approximately a straight line. Equation (6) then gives the familiar result $-t_o = t_R \simeq \frac{V_R}{c}$ where $\frac{V_R}{c} = \frac{1-\bar{\Omega}}{n} \simeq \frac{-\bar{\Omega}}{n}$ for $\bar{\omega} << 1$ and $E_R = \frac{1}{2}mV_R^2$ (Kennel and Petchek, 1966). Relativistic particles can interact with whistlers either by Doppler-shifting the wave frequency up to the local electron gyrofrequency or by lowering their gyrofrequency sufficiently by a mass increase, thus the resonance curve departs from a straight line. Induced emission or absorption of the whistler depends on the local derivatives of the distribution function along the resonance curve. The growth or damping of a mode depends on the net energy contribution to a wave from the particles. For a distribution such as $f \propto \sin^M\theta$, a given frequency may be unstable by equation (3) but for those particles with $p_z = 0$, $\left.\frac{\partial f}{\partial \theta}\right|_{\theta=\pi/2} = 0$ and consequently those particles are energized by the growing mode.

For application to energetic electrons in Jupiter's magnetosphere it is convenient to use the ultrarelativistic approximation where $t >> 1$. In that case the resonance curve can be approximated by $x \simeq \frac{1}{n}[1 - \bar{\Omega}/t]$ valid so long as $t_o \simeq \frac{\bar{\Omega}}{1+n} >> 1$. If the relativistic electron distribution obeys a power law in momenta (energy) we can model the distribution at the magnetic equator as $f(\vec{t}) = B\frac{\sin^M\theta}{t^{N+2}}$ where M and N are the pitch angle and spectral indices of the distribution and B is a normalization. A convenient normalization is in terms of the omnidirectional integral flux. For relativistic particles where $\gamma = \sqrt{1+t^2} \simeq t$ we have

$$f(\vec{t}) = \frac{(N-1)\gamma^{N-1} J(>\nu)}{4\pi c} \frac{2}{\sqrt{\pi}} \frac{\Gamma\left(\frac{M+3}{2}\right)}{\Gamma\left(\frac{M+2}{2}\right)} \frac{\sin^M\theta}{t^{N+2}} \qquad (7)$$

such that $J(>\gamma) = \int_\gamma^\infty d\gamma \int_{4\pi} d\Omega t^2 c f(\vec{t})$. Since the growth rate depends on the distribution above t_o by (3), if we consider frequencies such that t_o lies in the domain where (7) is valid, we do not have to specify the behavior of the distribution at low energies. When (7) is inserted into the growth rate and the change of variable $x = \frac{1}{n}[1 - \bar{\Omega}/t]$ is made we have the following expression for the logarithmic gain scaled to a Jovian radius at the magnetic equator:

$$k_i R_J = \frac{\pi^2 e R_J}{B_o(L)} \frac{J(>\gamma)}{c} \left[\frac{\gamma}{\bar{\Omega}/n}\right]^{N-1} F(N, M, u) \tag{8}$$

where

$$F(N, M, u) = (N - 1)\frac{1}{\sqrt{\pi}} \frac{\Gamma\left(\frac{M+3}{2}\right)}{\Gamma\left(\frac{M+2}{2}\right)} \int_{-1}^{u} dx(u - x)^{N-1} \cdot \left(1 - x^2\right)^{\frac{M}{2}} \left[u(N+2)\left(1 - x^2\right) + Mx(1 - ux)\right] \tag{9}$$

and $u = \frac{1}{n}$ is the normalized phase velocity of the wave. The growth rate in the relativistic approximation is dependent more fundamentally on the refractive index (actually phase velocity) rather than frequency.

If the quantity $t_R = \bar{\Omega}/n$ is considered as an effective resonant momentum, we can interpret equation (8) such that $\frac{\pi^2 e R_J}{B_o(L)}$ is the electromagnetic coupling of the wave to the distribution, $\frac{J(>\gamma)}{c}\left[\frac{\gamma}{\bar{\Omega}/n}\right]^{N-1}$ is the number "density" of particles with momenta $t \geq t_R$, and $F(N, M, u)$ is the kinematical resonance integral over the curvature of the distribution. Figure 2 is a plot of $F(N, M, u)$ for $N = 3$ and $M = 0, 2, 4, 6, 8$. Negative values of $k_i R_J$ correspond to growth. At low frequencies, $\bar{\omega} \to 0$, the growth rate goes to zero falling off like t_R^{1-N}. At higher frequencies where $u \sim 0.5$, the growth rate turns over and goes to zero since at a high enough phase velocity as much energy is taken from the wave as is given, the resonance curve in Fig. 1 becoming increasingly concave.

If the cold plasma density is known the growth rate can be computed as a function of frequency through the transformation $\frac{1}{u} = n(\omega)$. A basic role of the cold plasma density is to determine a range of possible values for the refractive index: $n_{MAX}(\omega = 0) \geq n(\omega) \geq n_{MIN}\left(\omega_p \text{ or } \frac{1}{2}\Omega_{ce}\right)$, n_{MIN} occurring at the lesser of ω_p or $\frac{1}{2}\Omega_{ce}$. This range specifies a phase velocity window which restricts the values of u in Fig. 2. For Jupiter, Frank et al. (1976) have reported cold proton densities, which have been used to plot Fig. 3. The cross-hatched region represents the accessible values of phase velocity at each $L = R/R_J$ on the magnetic equator. By comparison with Fig. 2, using representative values of $M = 4$ and $N = 3$ for the pitch angle and spectral indices, there exists a range of unstable frequencies for $1 \leq L \leq 12$.

In the non-relativistic regime one can determine the marginally stable frequency and the corresponding resonant particle energy above which whistlers are unstable. For relativistic particles, since the resonance curve is hyperbolic, the marginal stability point is model-dependent and a function of M and N. We can estimate $E_{M.S.}$ as follows. Choosing $N = 3$ and

$M = 4$, the marginally stable refractive $n_{M.S.} \sim 2$. If $n_{MIN} \gg 1$, $E_{M.S.}$ is non-relativistic and all relativistic electrons resonate with unstable whistlers. If $n_{MIN} = O(1)$, $t_o = t_{M.S.} \sim \frac{\bar{\Omega}_{M.S.}}{n_{M.S.}+1}$. Taking $\bar{\omega} \ll 1$ gives $t_{M.S.} = \frac{\Omega^2}{\omega_p^2}(n_{M.S.} - 1)$ and

$$E_{M.S.} \sim mc^2 \frac{\Omega^2}{\omega_{pc}^2}\left(n_{M.S.} - 1\right) \sim \frac{B^2}{4\pi n_o} \tag{10}$$

essentially the non-relativistic result.

b) Path Integrated Growth

The actual logarithmic power gain of a ducted whistler requires the path-integrated quantity $G = -2\int_{s_1}^{s_2} k_i(s)ds$ along a flux tube. With certain approximations the integration can be done analytically and the result is good at least to the accuracy that fluxes are measured. We treat the cold plasma density as constant along the flux tube and require equal contributions to the gain from above and below the magnetic equator. As the wave convects along the tube from one hemisphere to the other, the phase velocity changes as demonstrated in Fig. 2 by either of the arrows A or B depending on the phase velocity (frequency) at the initial point. Liemohn (1967) found that waves which were locally damped at the magnetic equator could have a net positive gain because of greater growth contributions at higher latitudes. This situation is demonstrated by the arrow A for the curve M = 4. However, the stably-trapped flux is determined by the waves with maximal gain, which case is represented by the arrow B; the situation marries the maximum of $F(M, N, u)$ with the maximum flux of resonant particles at the equator. But from inspection of Fig. 2, if we choose $N = 3$ and $M = 4$ the excursion due to $F(M, N, u)$ is small provided we consider phase velocities such that $u(\omega) \lesssim .35$. We thus approximate $F(M, N, u)$ as constant in this frequency range and having the constant value $F(N = 3, M = 4, u = 0) = -5/16$ (dotted lines) along the flux tube. The path integration then simply involves the scaling of the other factors in equation (8) along the flux tube. The coupling scales like (B_o/B) off the equator. For a $\sin^M\theta$ distribution the flux scales like $(B_o/B)^{M/2}$ (Roederer, 1970). If $n = 0(1)$, $t_R = \bar{\Omega}/n$ scales like (B/B_o). The growth rate along the tube is then approximately

$$k_i(s) = \left(\frac{B_o}{B}\right)^{\frac{M+2N}{2}} k_i(EQ) = \left(\frac{B_o}{B}\right)^{5} k_i(EQ) \qquad (11)$$

Then if $G = -4\int_o^{s_{max}} k_i(s)ds$ and using the harmonic approximation $B/B_o \simeq 1 + s^2/s_o^2$ for a dipole magnetic field where $s_o^2 = \frac{2}{9}L^2R_J^2$ we have

$$G \simeq -4k_i(EQ)\,LR_J\sqrt{\frac{2}{9}}\int_o^{y_{max}} dy\,(1+y^2)^{\frac{-(M+2N)}{2}} \qquad (12)$$

When $y_{max} = 1$ we are at the limit of the harmonic approximation, but since the integrand has decreased significantly, we extend the upper limit to infinity and for $N = 3$, $M = 4$ the integration yields $G = -4k_iR_JL(0.2)$. Thus

$$G(N = 3,\ M = 4) \simeq \frac{1}{4}\,\frac{\pi^2 e R_J L^4}{B_o}\,\frac{J(>\gamma)}{c}\left[\frac{\gamma}{\bar{\Omega}/n}\right]^2 \qquad (13)$$

We note that in scaling $t_R \simeq (\frac{B}{B_o})$ we required $n = 0(1)$. If the opposite extreme prevailed, $n \gg 1$, then $t_R \sim (\frac{B}{B_o})^{3/2}$ and $k_i(s) \simeq (\frac{B_o}{B})^6\, k_i(EQ)$, with the integration providing $G = -4k_i R_J L(0.18)$, a slight difference. All quantities in equation (13) refer to the equator and (13) is valid for frequencies such that $u(\omega) \lesssim 0.35$.

Generalizing this procedure for arbitrary values of M and N and approximating $F(N, M, u) \simeq F(N, M, u = 0)$ we have

$$G \simeq \sqrt{\frac{2}{9}}\,M(N-1)\,\frac{\Gamma(\frac{N+1}{2})\ \Gamma(\frac{M+3}{2})\ \Gamma(\frac{M+2N-1}{2})}{\Gamma(\frac{M+2N}{2})\ \Gamma(\frac{M+N+3}{2})}\,\frac{\pi^2 e R_J L^4}{B_o}\,\frac{J(>\gamma)}{c}\left[\frac{\gamma}{\bar{\Omega}/n}\right]^{N-1} \qquad (14)$$

We note that over the parameter range $M = 2, 4$, $N = 2, 3, 4$, G is a relatively weak function of the indices, viz. $G(N = 2, M = 4) \simeq (1.27)\ G(N = 4, M = 2)$ at the extreme. This is due to the compensating effects of, say, larger growth rate (M increases) with smaller effective path length from equation (11). Therefore, over this parameter range we will take equation (13) as valid generally.

c. Stably-Trapped Limit

Application of equation (13) to the concept of the stably-trapped limit (Kennel and Petchek, 1966) requires a knowledge of the power reflection coefficient, R, for the whistlers in the flux tube. If we treat the problem as strictly one-dimensional and define the volume emissivity η for parallel-propagating whistlers, the equation of radiative transfer $\frac{\partial I}{\partial s} = -2k_i(s)I + \eta(s)$ yields for the intensity

$$I(s_{MAX}) = \frac{\int_{-s_{MAX}}^{s_{MAX}} ds\,\eta(s)\,\exp\left[-2\int_{s}^{s_{MAX}} ds'\,k_i(s')\right]}{1 - Re^{G}} \quad (15)$$

It is clear that as $Re^G \rightarrow 1$, sufficiently high wave intensities will result that can relax the distribution on a time scale comparable to that of the particle source. No in situ wave measurements have been conducted at Jupiter, but if we assume that as in the Kennel-Petchek theory 5% of the wave energy is reflected, a gain of $G = \ln(1/R) = 3$ is sufficient to maintain the stably trapped equilibrium. Equation (13) then gives a stably-trapped flux limit of

$$J^*(>t_R) \simeq 4 \times 10^{10} L^{-4} \frac{\text{electrons}}{\text{cm}^2\ \text{sec}} \quad (16)$$

The non-relativistic Kennel-Petschek result scaled to Jupiter is

$$J^*_{N-R}(>E_R) = 7 \times 10^{10} L^{-4} \left(\frac{B_J R_E}{B_E R_J}\right) = 8.5 \times 10^{10} L^{-4} \quad (17)$$

which does not differ greatly from the relativistic result. This fact is borne out by inspection of (8). The relativistic aspect is manifest only in the resonance integral F(N, M, u) and if the distribution is not too different topologically in the non-relativistic regime, F(N, M, u) should not vary a great deal. This can be further motivated by assuming a power law distribution in non-relativistic energy. If we take the distribution with $\frac{1}{2}mv^2 = T$ as $f(\vec{v}) = B\frac{\sin^M\theta}{T^{N+1}}$ where M and N are the pitch angle and spectral indices, a non-relativistic analysis yields the result equivalent to equation (8) for $\bar{\omega} << 1$ of

$$k_i R_J = \frac{\pi^2 e R_J}{B_o} \frac{\left(T/T_R\right)^{N-1} J(>T)}{c} F_{N-R}(N, M, \bar{\omega}) \qquad (18)$$

where

$$T_R = \frac{1}{2}mV_R^2 \simeq \frac{1}{2}m\left(\frac{\bar{\Omega}c}{n}\right)^2$$

and

$$F_{N-R}(N, M, \bar{\omega}) = -M(N-1)\frac{2}{\sqrt{\pi}} \frac{\Gamma\left(\frac{M+3}{2}\right)}{\Gamma\left(\frac{M+2}{2}\right)}\left[1 - \frac{2\bar{\omega}}{M(1-\bar{\omega})}\right] \cdot \int_o^{\pi/2} d\theta \sin^{M+1}\theta \cos^{2N-1}\theta \qquad (19)$$

If $\bar{\omega} << 1$ and we neglect the second term in (19), then for N = 3 and M = 4, $F_{N-R} = -\frac{1}{4}$ compared to $F \simeq -\frac{5}{16}$ taken previously.

Equation (16) describes a limit which should not be exceeded statistically by a flux of electrons at any energy. If the flux $J(>\gamma)$ is known and obeys a power law, the distribution should depart from the power law and harden considerably at energies below the transition momentum given by

$$t_R = \gamma \left[\frac{J(>\gamma)}{J^*(>t_R)} \right]^{\frac{1}{N-1}} \tag{20}$$

since further extrapolation of the distribution by a power law to lower momenta would produce excessive amplification of the whistlers.

In Fig. 4 we have plotted $J^*(>t_R)$ for comparison with observation. For $L \leq 12$ measured electron spectra have spectral indices from 3 - 3.5 at high energies and pitch angle indices from 2 - 4 (Van Allen et al., 1974; Baker and Van Allen, 1976). We have taken the representative values of N = 3 and M = 4 throughout. The circles are equatorial electron fluxes of 5 Mev electrons from an empirical formula given by McIlwain and Fillius (1975); the diamonds are 5 Mev fluxes reported by Baker and Van Allen (1976) with error bars (D.N. Baker, private communication).

The theoretical uncertainty is more sensitive to the basic model we have chosen rather than to M and N values. We have taken the most efficient situation of parallel-propagation for which the growth rate is maximum (Kennel and Petschek, 1966); the finite $k_\perp$ effects should raise J^*. A reflection coefficient of 0.5% rather than 5% raises J^* by a factor of 1.8. Another possibility is that centrifugal forces at large L confine the cold plasma to within $\pm 1\ R_J$ of the equator as suggested by Ioannidis and Brice (1971). The effective path length is then just 2 R_J

and for N = 3, M = 4 the modification to the limiting flux becomes $J^* \sim 8 \times 10^{9} L^{-3} cm^{-2} sec^{-1}$ which is also plotted in Fig. 4. Within the experimental and theoretical uncertainties involved, the data lie near the stably-trapped limit and suggest that whistlers may be active and that the transition energy is approximately 5 Mev at low L values.

III DIFFUSION THEORY

a. Homogeneous Quasi-Linear Theory

If we consider a homogeneous plasma with parallel-propagating whistlers described by a one-dimensional electric field spectral density $\mathcal{E}(\omega)$ for waves travelling in just one direction, then

$$\langle \frac{E^2(\vec{x}, \tau)}{8\pi} \rangle = \mathcal{E}_{TOT} = 2\int_0^\infty d\omega \mathcal{E}(\omega);$$

the factor of two includes the waves travelling in the opposite sense. The quasilinear diffusion equation for the resonant electrons can be written

$$\left.\frac{\partial f(\vec{x}, \vec{t})}{\partial \tau}\right)_{Q.L.} = -\frac{\partial}{\partial \vec{t}} \cdot \vec{J} \tag{21}$$

where

$$\vec{J} = -\frac{4\pi^2 e^2}{m^2 c^2} \sum_{\omega_\pm} |V_g| \frac{\mathcal{E}(\omega)}{|V_g - v_z|} \left[\frac{\partial f}{\partial t_\perp} \pm \frac{n}{\gamma} \left[t_\perp \frac{\partial f}{\partial t_z} - t_z \frac{\partial f}{\partial t_\perp} \right] \right]$$
$$\left[\left(1 \mp \frac{n t_z}{\gamma} \right) \hat{t}_\perp \pm \frac{n t_\perp}{\gamma} \hat{t}_z \right] \tag{22}$$

and $\hat{t}_\perp$, $\hat{t}_z$ are unit vectors. Since a relativistic particle interacts with two waves going in opposite directions, the summation is over the frequencies satisfying $\omega_\pm \left[1 \mp n(\omega_\pm) \frac{v_z}{c} \right] = \frac{\Omega_{ce}}{\gamma}$ and $n(\omega_\pm) = \frac{|k| c}{\omega_\pm}$ for the R mode travelling in the $\pm$ direction along $\vec{B}_o$. The combination in the first bracket is proportional to the incremental growth rate (see equation (1)) and has a sign, $\sigma = +1$ (-1), corresponding to the particles giving (taking) energy from the wave. The diffusion current is thus proportional to the spectral density and incremental growth rate and, if we consider the interaction with just an upward travelling wave, has the direction given by the angle ψ_+ (see Fig. 5), where $\tan \psi_+ = \frac{J_z}{J_\perp} = \frac{nv \sin \theta}{c - nv \cos \theta}$. For $v \approx c$ we have

$$\psi_+ = \tan^{-1}\left[\frac{n \sin \theta}{1 - n \cos \theta}\right] + \frac{\pi}{2}(1 + \sigma) \qquad (23)$$

In general, ψ_+ is directed such that those particles giving energy increase their pitch angle, $\theta \to \pi$, and those receiving energy decrease their pitch angle, $\theta \to 0$. If $n \cos \theta << 1$, $\psi_+ \simeq \tan^{-1}[-\tan \theta] + \frac{\pi}{2}(1 + \sigma)$, which describes pitch angle diffusion along approximately iso-energy surfaces, since $t_z^2 + t_\perp^2$ = constant leads to $\frac{dt_z}{dt_\perp} = -\tan \theta$. At $\theta = \pi/2$ $\psi_+ = \tan^{-1} n$, but these particles react equally to downgoing waves and the net current is in the $\hat{t}_\perp$ direction. Thus 90° pitch angle particles should random-walk along the $t_\perp$ axis to higher energies. The features of the diffusion are sketched qualitatively in Fig. 5. The locus of points where the incremental growth is zero defines a cone inside of which the diffusion is generally directed outwards. We note that non-relativistically, if we use the resonance condition $1 - \frac{nv}{c} \cos \theta = \frac{1}{\gamma}\bar{\Omega} \simeq \bar{\Omega}$,

$$\psi_+ = \tan^{-1}[(\bar{\omega} - 1)\tan \theta] + \frac{\pi}{2}(1 + \sigma) \qquad (24)$$

and for low frequencies, $\bar{\omega} \ll 1$, pitch angle diffusion results.

b. Radial Diffusion Particle Source

The immediate source of relativistic electrons in the inner Jovian magnetosphere has been identified as inward radial diffusion resulting from third adiabatic invariation violation (see, e.g., Simpson et al., 1974). As a source term the process can be written

$$\left.\frac{\partial f(\vec{x}, \vec{t})}{\partial \tau}\right)_{R.D.} = L^2 \frac{\partial}{\partial L}\left[\frac{1}{L^2} D_{LL} \frac{\partial f}{\partial L}\right] \tag{25}$$

where the derivatives, $\left.\frac{\partial}{\partial L}\right)_{M,J.}$ are taken at constant first and second invariant, M and J. If we neglect synchrotron losses and bounce average, $\langle\rangle$, equation (21), the evolution of the distribution under combined radial diffusion and quasilinear momentum-space diffusion by parallel-propagating whistlers is described by

$$\frac{\partial f(\vec{x}, \vec{t})}{\partial \tau} = \left.\frac{\partial f}{\partial \tau}\right)_{R.D.} + \left\langle\left.\frac{\partial f}{\partial \tau}\right)_{Q.L.}\right\rangle \tag{26}$$

If electric field intensities are small the first term dominates, but if the distribution is sufficiently unstable to whistlers and in a stably-trapped equilibrium, quasilinear diffusion must occur on a time scale comparable to that of radial diffusion. The solution of (26) in the steady state is formidable, but we can extract a crude approximation of the bounce-averaged $\langle \mathcal{E}(\omega)\rangle$ in the stably-trapped regime. The radial diffusion coefficient can be modeled as $D_{LL} = D_o L^{\alpha}$ where $D_o \sim 10^{-10}\,\mathrm{sec}^{-1}$ and $\alpha \sim 4$ (Barbosa and Coroniti, 1975). Thus we estimate $\tau_{R.D.} \sim \frac{1}{D_o L^{\alpha-2}}$. Writing (21) in spherical coordinates, we estimate the Q.L. relaxation time as $\tau_{Q.L.} \sim \frac{m^2c^2t^2}{4\pi^2e^2}\frac{|V_g - v_z|}{\langle V_g \mathcal{E}(\omega)\rangle}$. The crudeness of the approximation is manifest in the neglect of the pitch angle dependence of $\tau_{Q.L.}$ and also by the neglect of the derivatives in (22) and (25) which to some degree are compensatory. Equating characteristic times yields

$$\langle V_g \mathcal{E}(\omega)\rangle = \langle I_E(\omega)\rangle \sim \frac{m^2c^2t^2}{4\pi^2e^2}|V_g - v_z| D_o L^{\alpha-2} \tag{27}$$

If $|V_g - v_z| \sim c$ and $t = t_R = \bar{\Omega}/n$, then letting $2\pi\nu = \omega$ we have for the magnetic field intensity of modes travelling parallel to the field in one direction

$$\langle I_B(\nu)\rangle = 2\pi\langle n^2 I_E(\omega)\rangle \sim 1.5 \times 10^{-18}\bar{\Omega}^2 L^{\alpha-2}\ \frac{\mathrm{WATTS}}{\mathrm{m^2 Hz}} \tag{28}$$

$\langle I_B\rangle$ is a lower limit such that intensities much lower than (28) will not produce significant decreases in phase space densities evolving under pure radial diffusion.

We can estimate typical fluctuation field strengths from (28). If the bandwidth $\Delta\nu \sim \nu \sim \frac{\Omega_{ce}}{2\pi}$, then $\frac{(\delta B)^2}{8\pi} \sim \frac{2}{V_g}\Delta\nu\langle I_B(\nu)\rangle$ and $\delta B \sim .5\ m\gamma L^{\frac{1}{2}(\alpha-5)} = .5\ m\gamma L^{-\frac{1}{2}}$ for $\alpha = 4$. Such fields should be detectable by future spacecraft to Jupiter.

IV DISCUSSION

The evaluation of the limiting flux assumed whistlers were generated locally and did not propagate across L-shells. Without more detailed knowledge of the cold plasma distribution, ray path computations would be speculative. However, the fact that observed fluxes are lower than the limiting flux inside of $L \leqslant 4$, where phase space losses are still apparent, suggests that whistlers may propagate inward.

ACKNOWLEDGEMENT

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FIGURE CAPTIONS

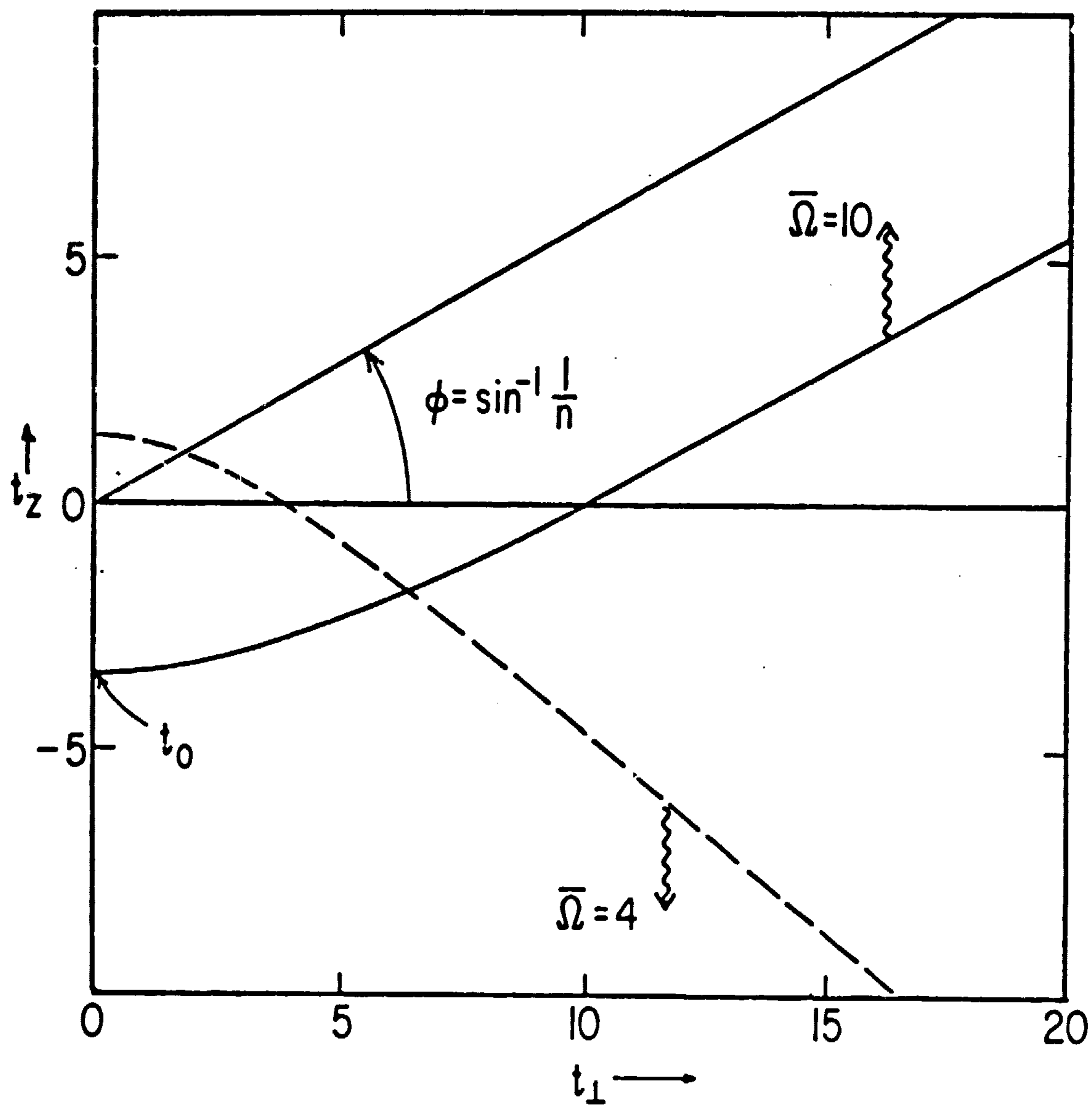
Figure 1. Resonance curves plotted in normalized momentum space for $R = \frac{\omega_p^2}{\Omega_{ce}^2} = \frac{1}{4}$. The solid line is the resonance with a whistler ($\bar{\Omega} = 10$, $n = 1.94$) travelling along the field; the dashed line with a whistler ($\bar{\Omega} = 4$, $n = 1.53$) travelling anti-parallel.

Figure 2. Plot of the resonance function F(M, N, u) for M = 0, 2, 4, 6, 8 and N = 3. Growth occurs for negative values of F(M, N, u) increasing with M. The arrows represent convective changes of F along the flux tube.

Figure 3. Accessible phase velocities for whistlers using Frank et al. (1976) cold plasma observations.

Figure 4. Plot of the limiting flux $J^*(>t_R)$. The circles are 5 Mev equatorial fluxes of McIlwain and Fillius (1975); the diamonds are 5 Mev fluxes of Baker and Van Allen (1976) with error bars. The dotted line is the modification to J* from centrifugal effects.

Figure 5. Qualitative view of diffusion in normalized momentum space arising from interaction with whistler propagating parallel to $\vec{B}_o$ only. The small arrows give only the direction of the diffusion. The cone edge is defined by the integrand of (3) being zero for N = 3 and M = 4.



5
0
-5
0
5
10
15
20
t_z
$t_\perp$
$\phi = \sin^{-1}\frac{1}{n}$
$\bar{\Omega}=10$
$\bar{\Omega}=4$
t_0

FIG. 1

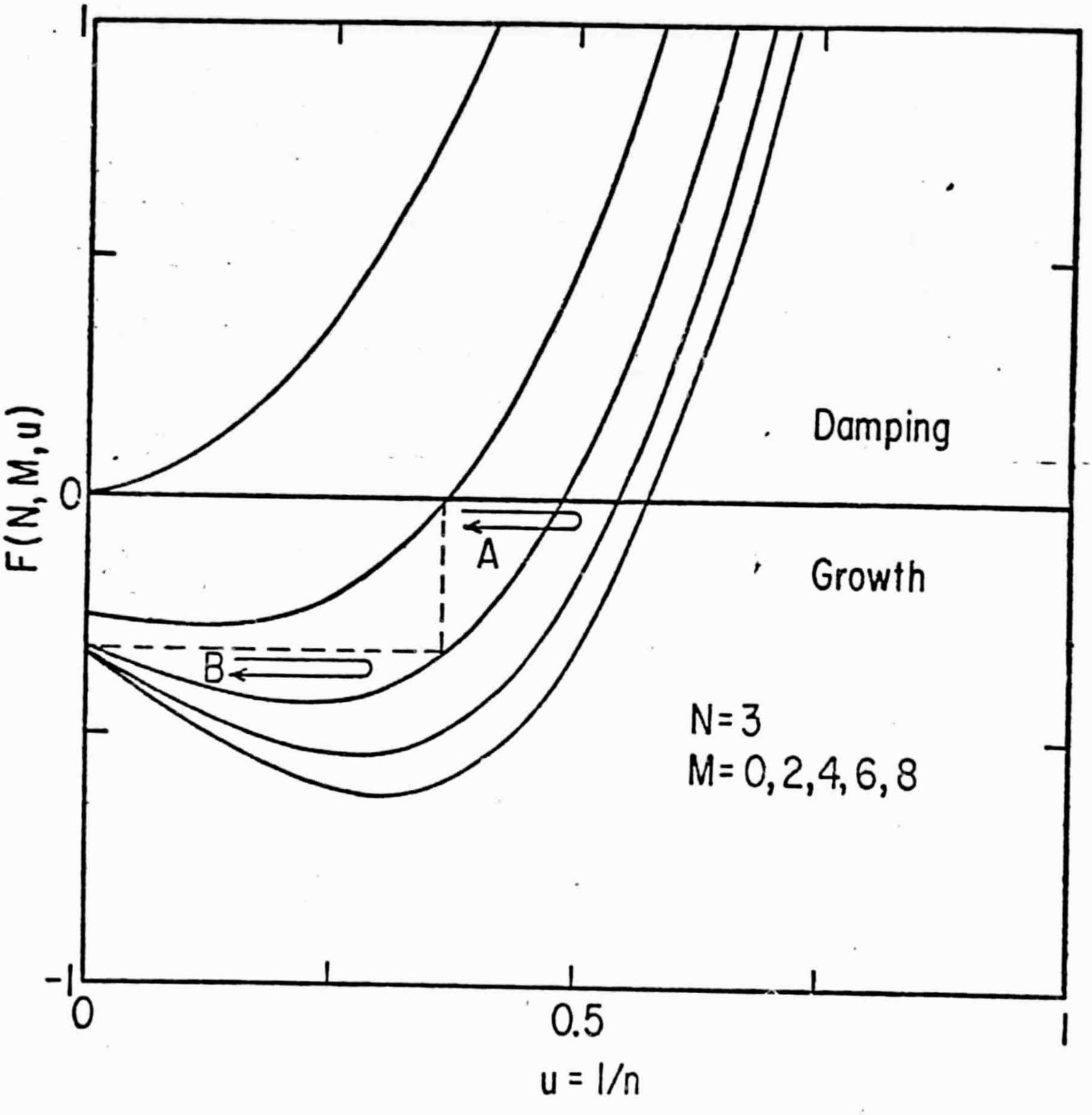


FIG. 2

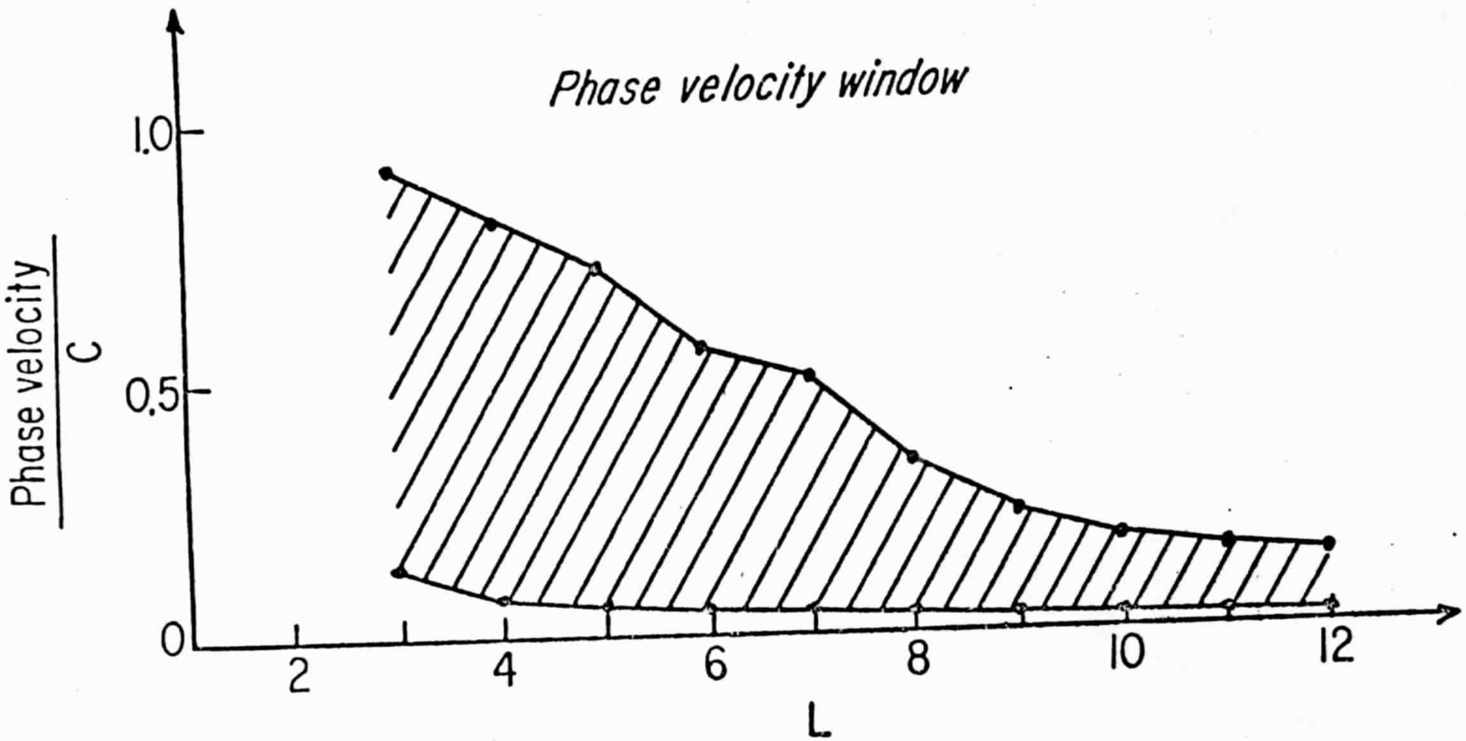
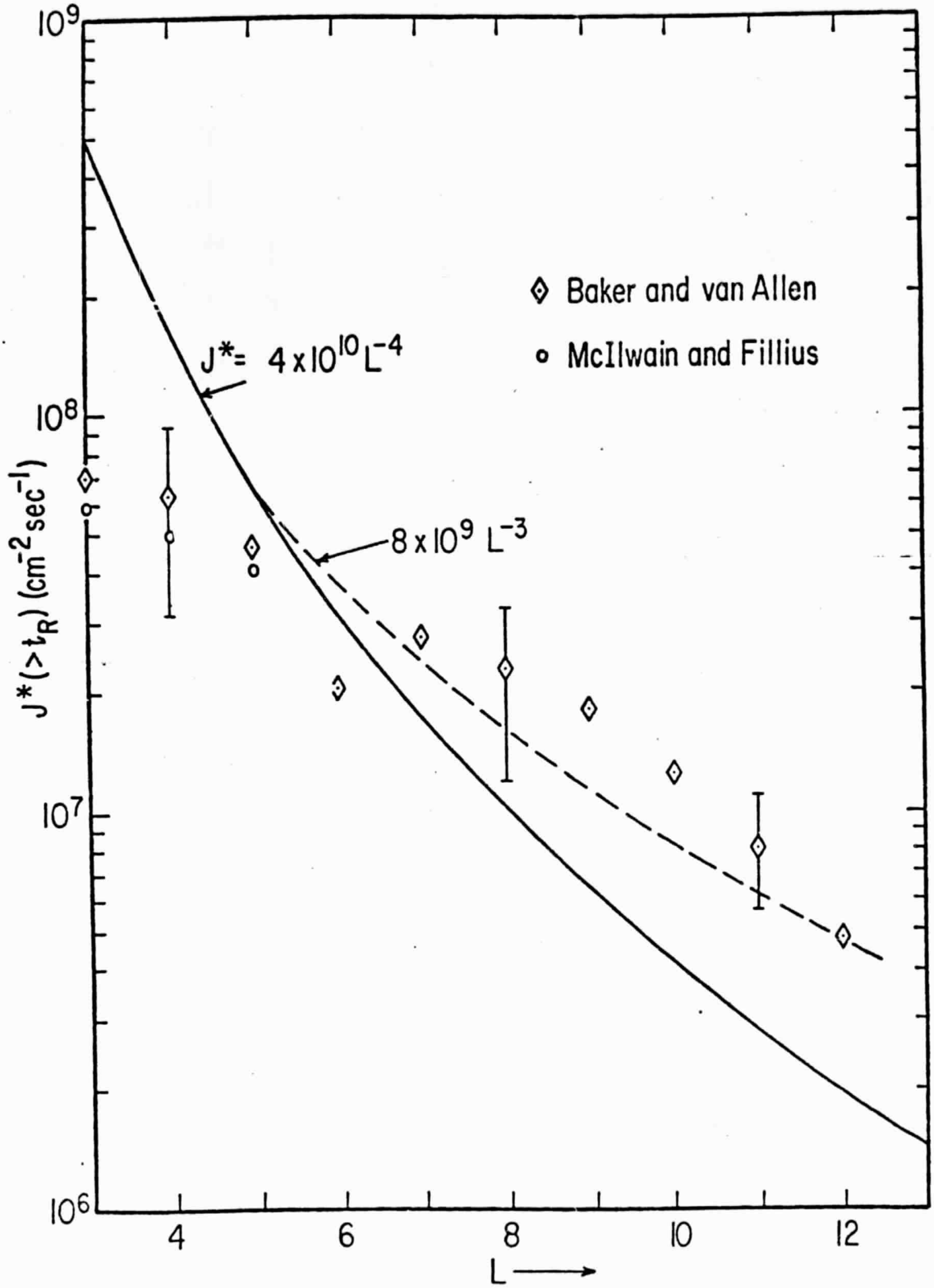


FIG. 3



10^9
10^8
10^7
10^6
$J^*(>t_R)$ $(cm^{-2} sec^{-1})$
Baker and van Allen
McIlwain and Fillius
$J^* = 4 \times 10^{10} L^{-4}$
$8 \times 10^9 L^{-3}$
4
6
8
10
12
L

FIG. 4

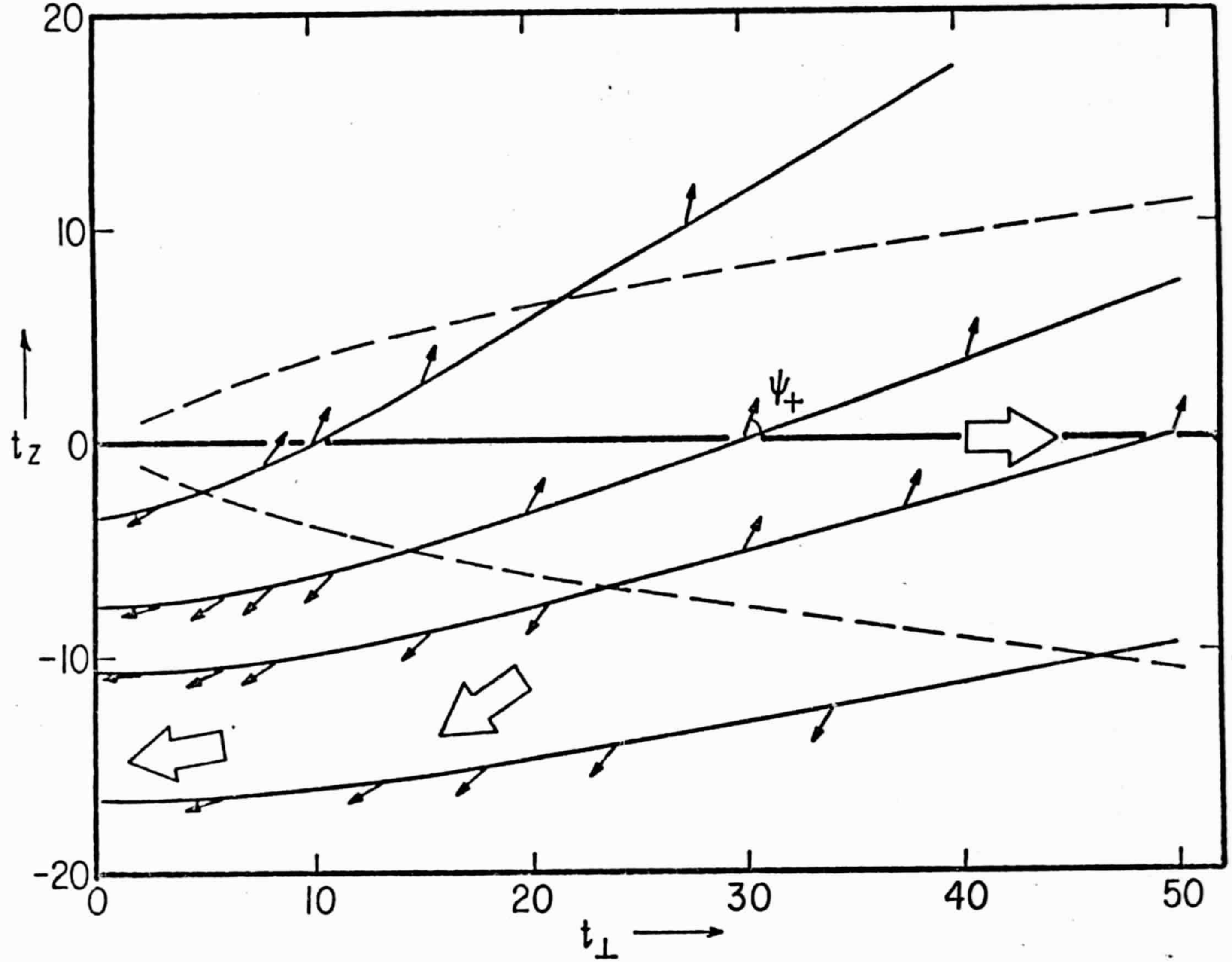


FIG 5

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