If the level in question had negative parity, it would be formed by p capture and would de-excite to the ground state by $E1$ radiation. If its parity were positive, it would be formed by s capture and de-excite to the ground state by (pure) $M1$ radiation. It is feasible to make estimates of the expected proton and radiation widths for both possibilities, and the results of both of these estimates are favorable to an assignment of positive parity.

In the case of Γ_p, the expected value is

$$\Gamma_p = (2k/\pi K)v_l D,$$

where the nomenclature adopted is given in I. A $\frac{1}{2}^-$ assignment to the resonance level would require a value of 11 Mev for D. A $\frac{1}{2}^+$ assignment on the other hand would give $D=2.6$ Mev. In a recent compilation of proton widths by C. van der Leun, of this laboratory, an average value of D equal to 2 Mev was found for 6 levels in this region of the periodic table. The value of 2.6 Mev corresponding to a $\frac{1}{2}^+$ assignment is closer to the average than that corresponding to a $\frac{1}{2}^-$ assignment. This argument is not considered to be conclusive, but does render the $\frac{1}{2}^+$ value the most probable.

Our measured value of $\Gamma_{\gamma 0}$ also points to a $\frac{1}{2}^+$ assignment. According to the latest estimate of Wilkinson,[5] based on a study of dipole radiation widths for $A \leqslant 20$, widths of 16 ev and 1.1 ev would be expected for $E1$ and $M1$ transitions of this energy, respectively. The latter, corresponding to a $\frac{1}{2}^+$ assignment, is seen to be in agreement with all measured values.

On the basis of the above discussion it is felt to be reasonably certain that the level in question is of spin and parity $\frac{1}{2}^+$.

ACKNOWLEDGMENTS

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Our thanks are due our colleagues A. M. Hoogenboom, E. Gelsema, C. van der Leun, and Dr. T. Muller for invaluable help in the setting up and execution of this experiment, and in the calculations.

[5] D. H. Wilkinson, author in *Proceedings of the Rehovoth Conference*, edited by H. J. Lipkin (North-Holland Publishing Company, Amsterdam, 1958), p. 175.

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Cosmic-Ray Modulation by Solar Wind*

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It is shown that the hydrodynamic outflow of gas from the sun observed by Biermann results in a reduction of the cosmic-ray intensity in the inner solar system during the years of solar activity. The computed cosmic-ray energy spectrum so closely resembles the observed spectrum at earth that we suggest the outflow of gas to be the explanation for the 11-year variation of the cosmic-ray intensity.

It is also suggested that perhaps the Forbush-type decrease, which is a local geocentric phenomenon, is the result of disordering of the outer geomagnetic field by the outflowing gas from the sun.

I. INTRODUCTION

IN a recent paper[1] it was shown that the hydrodynamic flow[2] of gas outward in all directions from the sun, as observed by Biermann,[3] stretches out the magnetic lines of force of the solar magnetic fields, and leads to an essentially radial magnetic field in the inner solar system; the field density at the point (r,θ,ϕ), well removed from the sun is

$$B(r,\theta,\phi)\cong B(a,\theta,\phi)(a/r)^2,$$

where a is the radius (7×10^5 km) of the sun, r is distance from the center of the sun, and θ and ϕ are the usual polar and azimuthal angles. A field density of 1 gauss at the sun[4] yields about 2×10^{-5} gauss at the orbit of earth ($r=1.5\times10^8$ km).

The high velocity (500–1500 km/sec) and low density ($\sim$500 ions/cm^3) of this outward streaming gas, the solar wind, leads to anisotropic thermal motions as a consequence of the anisotropic expansion. And it was pointed out[1] that when the gas pressure in the direction parallel to the magnetic field exceeded the pressure perpendicular by an amount Δp which was greater than $B^2/4\pi$, the dynamical equations[5] for the plasma predict instability of the magnetic field to transverse

* Assisted in part by the Office of Scientific Research and the Geophysics Research Directorate, Air Force Cambridge Research Center, Air Research and Development Command, U. S. Air Force.
[1] E. N. Parker, Phys. Rev. **109**, 1874 (1958).
[2] E. N. Parker, Astrophys. J. (to be published).
[3] L. Biermann, Z. Astrophys. **29**, 274 (1951); Z. Naturforsch. **7a**, 127 (1952); Observatory **107**, 109 (1957).
[4] H. W. Babcock and H. D. Babcock, Astrophys. J. **121**, 349 (1955).
[5] E. N. Parker, Phys. Rev. **107**, 924 (1957).

perturbations. It was estimated that the instability should set in somewhere in the vicinity of the orbit of earth, thereby leading us to the conclusion that the inner solar system is surrounded by a thick shell of disordered magnetic field of perhaps 2×10^{-5} gauss.

The existence of an enclosing shell of disordered field of $\sim10^{-5}$ gauss had been inferred earlier from considerations[6] of the decay of the cosmic-ray intensity from the solar flare of February 22, 1956, which suggested that the shell had (at that time) a thickness of about 4 astronomical units, and extended from about the orbit of Mars to the orbit of Jupiter.

Such a heliocentric shell of disordered field will, in the presence of the 500–1500 km/sec solar wind, have significant effects on the cosmic-ray intensity in the inner solar system.[7] Each irregularity of the magnetic field in the disordered shell is carried outward in the solar wind. The cosmic rays will tend to be swept ahead of the field fluctuations. Equilibrium is reached when the rate at which cosmic rays can diffuse in through the shell is equal to the rate at which they are removed by convection. Hence the cosmic-ray intensity inside the shell is less than outside.

In addition to this heliocentric reduction of the cosmic-ray intensity there may be a further effect of the solar wind on the cosmic-ray intensity. Polar magnetic stations show that the magnetic lines of force extending beyond about 6 earth's radii (those lines of force which come down within about 25° of the magnetic pole) are continually agitated by the solar wind. And we may infer from this that the lines of force are always slightly disordered beyond 6 radii; and when the solar wind is high, the disordering is probably substantial. Such disordering of the outer geomagnetic field affects the entrance of cosmic rays, and thereby causes an observer on the surface of earth to see a geocentric change in intensity.

II. HELIOCENTRIC MODULATION

It has been pointed out elsewhere[1] that we can expect disordering of the magnetic field in the heliocentric shell with scales up to a maximum of about $l=2\times10^6$ km; the proportionately slower growth of larger scales does not allow their full development in the time available.

The radius of curvature $P(\eta)$ of a cosmic-ray particle with rest energy Mc^2 and kinetic energy ηMc^2 is

$$P(\eta)=Mc^2[\eta(\eta+2)]^{\frac{1}{2}}/ZeB. \tag{1}$$

Hence the path of a 1-Bev proton in the shell field of $\sim2\times10^{-5}$ gauss has a radius of curvature of about 2.6×10^6 km, which we note is larger than the scale l of the disordering of the field in the heliocentric shell. Thus upon passing through the shell, a cosmic-ray proton undergoes a large number of random angular deflections of the order of $l/P(\eta)$. The largest deflection expected for a 1-Bev proton is, accordingly, of the order of 2/2.6 radians or about 45°; the smaller-scale ($<2\times10^6$ km) magnetic irregularities, and higher-energy particles, will result in smaller deflections. But the cumulative effect of n small angular deflections l/P is

$$\Theta(\eta)\cong n^{\frac{1}{2}}l/P(\eta).$$

Therefore, if we define a collision to be a total deflection of $\Theta(\eta)=\pi/2$, then the mean free path is

$$L(\eta)=\pi^2P^2(\eta)/4l \tag{2}$$

for a particle with energy η.

Elementary kinetic theory tells us that the coefficient of diffusion of particles with a velocity w and mean free path L is

$$\begin{aligned}\kappa(\eta)&\cong\tfrac{1}{3}w(\eta)L(\eta)\ \text{cm}^2/\text{sec}\\&=\frac{\pi^2}{12}\left(\frac{M^2c^5}{Z^2e^2B^2l}\right)\frac{[\eta(\eta+2)]^{\frac{3}{2}}}{(\eta+1)}.\end{aligned} \tag{3}$$

If the diffusing medium (the magnetic irregularities) has a general motion $\mathbf{v}$, the diffusive and convective transport equation is[7]

$$\partial j(\eta)/\partial t=-\nabla\cdot[\mathbf{v}j(\eta)]+\nabla\cdot[\kappa(\eta)\nabla j(\eta)] \tag{4}$$

for the density $j(\eta)$ of cosmic-ray particles with energy η.

Supposing that the disordered shell extends uniformly, and with spherical symmetry, from a solar distance $r=r_1$ out to $r=r_2$, we have upon integration of (4) that the steady-state cosmic-ray density $j_0(\eta)$ inside the shell is related to the galactic density $j_\infty(\eta)$ outside by

$$j_0(\eta)=j_\infty(\eta)\exp\left\{-\frac{12v(r_2-r_1)lZ^2e^2B^2(\eta+1)}{\pi^2M^2c^5[\eta(\eta+2)]^{\frac{3}{2}}}\right\}. \tag{5}$$

If we suppose that we have the typical values mentioned above, $l\cong2\times10^6$, $B\cong2\times10^{-5}$ gauss, $r_2-r_1=4$ astronomical units (6×10^{13} cm), and $v=10^3$ km/sec, then we have

$$j_0(\eta)=j_\infty(\eta)\exp\{-2.16(\eta+1)/[\eta(\eta+2)]^{\frac{3}{2}}\} \tag{6}$$

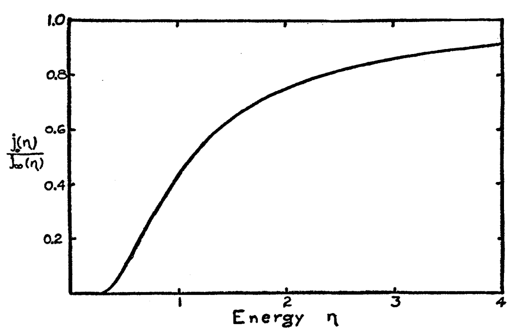

FIG. 1. Heliocentric depression of the cosmic-ray intensity in the inner solar system by the solar wind in the heliocentric shell of disordered magnetic field.

[6] Meyer, Parker, and Simpson, Phys. Rev. **104**, 768 (1956).
[7] E. N. Parker, Phys. Rev. **103**, 1518 (1956).

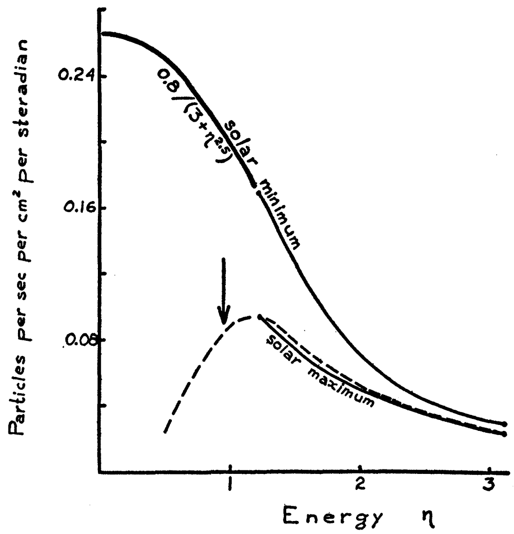

FIG. 2. The solid curves represent the observed primary cosmic-ray differential energy spectrum during the 1954 period of minimum solar activity and during the years of maximum activity; the curve $0.8/(3+\eta^{2.5})$ below 1.2 Bev is a qualitative representation of the low-energy observations; the vertical arrow indicates the energy at which the spectrum was observed to drop off rapidly during solar activity. The broken line is the spectrum which results from the solar wind in the heliocentric shell.

for protons. $j_0(\eta)/j_\infty(\eta)$ is plotted in Fig. 1. Note that the result is a depression of the cosmic-ray spectrum at all energies; but particularly at low energies, where a sharp cutoff occurs. The mean particle energy inside the shell is increased as a consequence of the decreased intensity; the low-energy cosmic-ray particles are kept out of the inner solar system by the solar wind in the heliocentric shell.

But note that this is just the sort of depression of the cosmic-ray intensity that one observes during the 11-year cycle of solar activity. When solar activity is at a minimum (as it was in 1954), the differential energy spectrum of the primary cosmic rays is observed[8] to be approximately of the form $j(\eta)\propto\eta^{-n}$, with $n\cong 2.7$ above about 2 Bev/nucleon; below 2 Bev/nucleon n has a somewhat smaller value, but there is no indication[9] that $j(\eta)$ reaches a maximum before zero energy.

On the other hand, with the reappearance of solar activity following the minimum, the cosmic-ray intensity begins to decrease. Meyer and Simpson[8] have shown that the cosmic-ray intensity at 2 Bev was depressed by 50% during 1948 when the sun was fairly active. Meredith, Van Allen, and Gottlieb[10] found no evidence for particles below about 1 Bev/nucleon in 1952 (the sun was still active), indicating a rather sharp low-energy cutoff. The depression of the cosmic-ray intensity during the years of solar activity is obviously due to the disappearance of the low-energy particles from the spectrum, as is shown directly by the shift of the latitude knee toward the equator; the mean energy of the particles observed at earth increases with decreasing cosmic-ray intensity.

We exhibit the observed cosmic-ray differential energy spectrum in Fig. 2 during the years of minimum and maximum solar activity. We have used the data of Meyer and Simpson[8] above 1.2 Bev for the solar minimum of 1954 and for the preceding solar maximum. We have included, in a qualitative way, Neher's observation[9] of the high abundance of low-energy particles during the solar minimum of 1954 by supposing that the spectrum is something like $0.8/(3+\eta^{2.5})$ particles per sec per cm² per steradian below 1.2 Bev. The vertical arrow at 0.9 Bev represents the energy below which Van Allen *et al.*,[10] found a rapidly diminishing particle density during solar maximum.

Now, suppose that during the solar minimum of 1954 the solar wind was so low that the observed spectrum was very nearly the galactic cosmic-ray spectrum (always to be found outside the solar system). Then it follows that if we should multiply the 1954 spectrum by the expected depression factor $j_0(\eta)/j_\infty(\eta)$ given in (6), we would obtain some rough idea of what spectrum we might expect to observe in the inner solar system during the years of solar activity, at least so far as the effects of the heliocentric shell are concerned. The broken line in Fig. 2 represents this expected spectrum; it agrees in its general characteristics with the observed spectrum during solar activity. The difference between the observed spectrum and the spectrum produced by the solar wind is so slight that we conclude the 11-year cosmic-ray variation to be principally the result of the solar wind in the heliocentric shell of disordered magnetic field.

III. GEOCENTRIC MODULATION

It has been pointed out elsewhere[7] that the onset of a Forbush-type decrease in the observed cosmic-ray intensity is sometimes so abrupt and geographically irregular that it must be a local geocentric effect.[11] We do not propose here to establish a final model for bringing about a geocentric depression of the cosmic-ray intensity. But we shall point out at least one of the more obvious possibilities for its occurrence. Then perhaps when we understand the dynamical interaction of the solar wind with earth in a little more detail, we shall be able to draw a more complete picture.

It was mentioned earlier that magnetic observing stations within about 25° of the geomagnetic poles observe a continued agitation of the field, which is evidently caused by the agitation of the outer regions of the geomagnetic field by the solar wind. A latitude

[8] P. Meyer and J. A. Simpson, Phys. Rev. **99**, 1517 (1955).
[9] H. V. Neher, Phys. Rev. **103**, 228 (1956); **107**, 588 (1957).
[10] Meredith, Van Allen, and Gottlieb, Phys. Rev. **99**, 198 (1955).

[11] Though, in view of the solar wind and radial solar field, we can no longer seriously entertain the suggestion made at that time[7] that the Forbush decrease is the result of terrestrial gravitational capture of passing interplanetary gas clouds containing disordered magnetic fields. Nor can we suppose any longer that the 11-year cycle is simply an accumulation of Forbush decreases during the years of solar activity.

of 65° (25° from the pole) corresponds to magnetic lines of force which extend outward to about 6 earth's radii, indicating the nearness of approach of the disturbing solar wind.

Suppose that the solar wind can perturb the geomagnetic field to within a radial distance $R=R_1$ at which point the magnetic energy density is equal to the kinetic energy density $\frac{1}{2}NMv^2$ of the solar wind. Then in the equatorial plane

$$R_1=b(B_0^2/4\pi NMv^2)^{1/6}, \tag{7}$$

where b is the radius of earth (6.4×10^8 cm) and B_0 is the horizontal component (about 0.4 gauss) of the geomagnetic field at the equator. Even on the quietest days, when $N\cong10^2$ ions/cm^3 and $v\cong500$ km/sec, we have $R_1=5.6b$, in agreement with the magnetic observations. And when the solar wind is high, with $N\cong10^5$ ions/cm^3 and $v\cong1500$ km/sec, we have $R_1=1.24b$.

Thus the geomagnetic field can be penetrated to a considerable depth by tongues of ionized gas from the solar wind, and we conclude that there is no reason to believe that beyond R_1 the field bears any resemblance to a dipole: Besides being disordered, the outer field may be stretched out like smoke from a chimney in a high wind.

We expect that such disordering will affect the entrance of cosmic-ray particles to earth. Unfortunately, we do not have even a qualitative notion of the nature of the disordering; and it follows that we cannot deduce the quantitative effect on the cosmic-ray intensity.

However, we need a point of departure for future interpretation of the observational data and for theoretical investigation. Therefore, with the understanding that our results need not bear any significant relation to the actual case, consider entrance of cosmic-ray particles to the surface of earth at the geomagnetic pole with the assumption that beyond some distance $R=R_2$ (R_2 is not necessarily equal to R_1) the geomagnetic field is completely disordered, and that inside R_2 the field is the conventional dipole.

We shall restrict ourselves to particles (protons) with energies rather less than 10 Bev so that we may use the adiabatic invariant

$$\sin^2\theta/B\cong\text{constant}, \tag{8}$$

where θ is the angle of pitch, *viz.* the angle between the particle velocity $\mathbf{w}$ and the magnetic field, of density B.

It follows from (8) that a particle which has an angle of pitch θ where the field density is B cannot penetrate to where the field density is in excess of $B/\sin^2\theta$. Thus a cosmic-ray particle starting far out in the geomagnetic field, where B is small, must have an exceedingly small angle of pitch if it is to reach the surface of earth at $R=b$. The angle of pitch at $R=R_2$ must not exceed θ_c, where

$$\sin\theta_c=(b/R_2)^{\frac{3}{2}}. \tag{9}$$

We shall let $f(\eta,\theta,z)d\theta d\eta$ represent the number of cosmic-ray particles per unit volume, a distance z beyond R_2 in the disordered magnetic field, and which have energies in $(\eta,\eta+d\eta)$ and velocity $\mathbf{w}$ lying in the angular interval $(\theta,\theta+d\theta)$ from the radial or z direction. We shall suppose that the disordered field stretches uniformly from $R=R_2$ (at which point we put $z=0$) to $R=R_2+a$ (where $z=a$). Then if a is rather less than R_2 we have essentially a one-dimensional diffusion of cosmic rays from $z=a$ into the undistorted geomagnetic field beginning at $z=a$.

We represent the diffusion coefficient by $\kappa(\eta)$ and suppose that on the average a particle of energy η is isotropically scattered after a mean life $\tau(\eta)$. It follows that

$$\frac{\partial f(\eta,\theta,z)}{\partial t}=\kappa(\eta)\frac{\partial^2 f(\eta,\theta,z)}{\partial z^2}-\frac{f(\eta,\theta,z)}{\tau(\eta)}+\frac{\sin\theta}{2\tau(\eta)}\int_0^\pi du f(\eta,u,z). \tag{10}$$

We seek steady-state solutions of (10) subject to the boundary conditions that: (a) at the outer surface $z=a$ of the disordered field, $f(\eta,\theta,z)$ is just equal to the isotropic cosmic-ray distribution to be found throughout the inner solar system

$$f(\eta,\theta,a)=\tfrac{1}{2}j_0(\eta)\sin\theta; \tag{11}$$

(b) the particles at $z=0$ with $\theta>\theta_c$ are reflected by the increasing geomagnetic field back into the disordered field so that there is no net flux of particles across $z=0$, and hence at $z=0$

$$\kappa(\eta)\partial f(\eta,\theta,z)/\partial z=0 \tag{12}$$

for $\theta>\theta_c$; (c) the particles at $z=0$ with $\theta<\theta_c$ move freely into earth where they are absorbed, with the result that the flux at $z=0$ is

$$\kappa(\eta)\partial f(\eta,\theta,z)/\partial z\cong w(\eta)f(\eta,\theta,z) \tag{13}$$

for $\theta<\theta_c$. We have represented the particle velocity by $w(\eta)$.

The reader will readily verify that the steady-state solution of (10) which satisfies the boundary conditions (11), (12), and (13), is

$$f(\eta,\theta,z)=\tfrac{1}{2}j_0(\eta)\sin\theta\left\{1+C\left[\frac{z-a}{h(\eta)}\cos\frac{a}{h(\eta)}(1-\cos\theta_c)+q(\eta,\theta)\sinh\frac{z-a}{h(\eta)}\right]\right\}, \tag{14}$$

where

$$h^2(\eta)=\kappa(\eta)\tau(\eta), \tag{15}$$

$$C=w(\eta)/\kappa(\eta)Q(\eta,\theta_c), \tag{16}$$

$$Q(\eta,\theta_c)=\left\{\frac{1}{h(\eta)}\left[1+\frac{aw(\eta)}{\kappa(\eta)}\right](1-\cos\theta_c)+\left[\frac{w(\eta)}{\kappa(\eta)}\sinh\frac{a}{h(\eta)}+\frac{1}{h(\eta)}\cosh\frac{a}{h(\eta)}\right](1+\cos\theta_c)\right\}, \quad (17)$$

and

$$q(\eta,\theta)=\begin{cases}-(1-\cos\theta) & \text{if } \theta>\theta_c\\ +(1+\cos\theta) & \text{if } \theta<\theta_c.\end{cases} \quad (18)$$

Since θ_c is small, it follows that the particle density in $\theta<\theta_c$ at $z=0$, which makes up the cosmic-ray flux observed at earth, is reduced from the usual $\frac{1}{2}j_0(\eta)\sin\theta$ to

$$f(\eta,\theta,0)\cong\tfrac{1}{2}j_0(\eta)\sin\theta\left\{1-\frac{[w(\eta)/\kappa(\eta)]\sinh[a/h(\eta)]}{[w(\eta)/\kappa(\eta)]\sinh[a/h(\eta)]+[1/h(\eta)]\cosh[a/h(\eta)]}+O^2(\theta_c)\right\}.$$

Now if we suppose, after the manner of elementary kinetic theory, that the diffusion coefficient $\kappa(\eta)$ is related to the mean free path, $L(\eta)=w(\eta)\tau(\eta)$ by the expression

$$\kappa(\eta)=\tfrac{1}{3}L(\eta)w(\eta),$$

then it follows that $3h^2(\eta)=L^2(\eta)$ and $w(\eta)/\kappa(\eta)=3/L(\eta)$. Hence

$$f(\eta,\theta,0)\cong\tfrac{1}{2}j_0(\eta)\sin\theta\{1-[1+3^{-\frac{1}{2}}\coth(a3^{\frac{1}{2}}/L(\eta))]^{-1}+O^2(\theta_c)\}.$$

Since we are interested in cases where $f(\eta,\theta,0)$ is only 10 or 20% less than $\frac{1}{2}j_0(\eta)\sin\theta$,[12] it follows that $a/L(\eta)\ll 1$ and

$$f(\eta,\theta,0)\cong\tfrac{1}{2}j_0(\eta)\sin\theta L(\eta)/[3a+L(\eta)].$$

If we suppose that the field beyond R_2 is disordered on on a scale which is small compared to the radius of curvature of the trajectory of a particle in the field, then $L(\eta)$ is given by (2), and

$$f(\eta,\theta,0)\cong\tfrac{1}{2}j_0(\eta)\sin\theta\left[1+\frac{12alZ^2e^2B^2}{\pi^2M^2c^5\eta(\eta+2)}\right]^{-1}. \quad (19)$$

On the other hand, if the scale of the disordering is not small compared to the radius of curvature, we have $L(\eta)\cong P(\eta)$, and

$$f(\eta,\theta,0)\cong\tfrac{1}{2}j_0(\eta)\sin\theta\left[1+\frac{3aZeB}{Mc^2[\eta(\eta+2)]^{\frac{1}{2}}}\right]^{-1}. \quad (20)$$

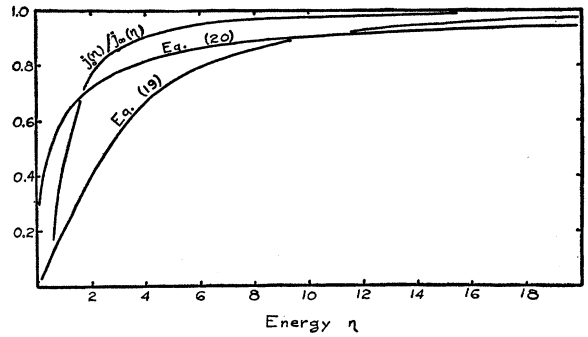

FIG. 3. The geocentric depression $f(\eta,\theta,0)/[\frac{1}{2}j_0(\eta)\sin\theta]$ of the cosmic-ray intensity at the magnetic poles by the disordering of the outer geomagnetic field by the solar wind. Equation (19) assumes that the disordering is on a small scale, and Eq. (20) on a large scale. The spectrum of the heliocentric depression $j_0(\eta)/j_\infty(\eta)$ is plotted for comparison.

The reduction, given in (19) and (20), of the cosmic-ray intensity observed at earth, $f(\eta,\theta,0)/\frac{1}{2}j_0(\eta)\sin\theta$, is shown in Fig. 3 for the case in which the intensity is depressed by 10% at 10 Bev. The form of the heliocentric reduction $j_0(\eta)/j_\infty(\eta)$ is also shown for purposes of comparison. Note how much flatter is the spectrum of the depression by the geomagnetic disordering, particularly (20).

We do not have any notion as to the expected scale of the small irregularities assumed in (19), so we cannot evaluate the consequences of our assumed 10% reduction. Suppose, however, that the thickness of the region of disordered field is $a\cong 5b\cong 3\times10^4$ km. Then a 10% reduction of the incoming 10-Bev proton intensity is achieved in (20) if $R_2\cong 10b$ (6.4×10^4 km), where the geomagnetic field density is 4×10^{-4} gauss. This seems a modest requirement inasmuch as we have estimated that the solar wind can readily push to $R\cong 6b$ (4×10^4 km). Consequently, we suggest that it may be the disordering of the outer geomagnetic field by an enhanced solar wind which is responsible for the onset of a Forbush decrease.

Unfortunately, we cannot treat the expected decrease of cosmic-ray intensity near the equator as we have done at the poles, without having a more detailed picture of the disordering. Therefore we cannot at the present predict how the Forbush decrease should vary with geomagnetic latitude.

[12] Note that in the limit of an infinitely thick and dense disordered field $[a/L(\eta)\gg 1]$, we have $f(\eta,\theta,0)\cong\frac{1}{2}j_0(\eta)\sin\theta/(1+\sqrt{3})$, which is a reduction to 0.37 of the density in the absence of disordered field.