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PLASMA IN ASTROPHYSICS

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Abstract

Two examples in which plasma theory is essential to the understanding of astrophysical phenomena are presented.

I. Introduction

In this review I should like to discuss two examples of plasma phenomena of importance to astrophysics. These are examples where astrophysical understanding hinges on further progress in plasma physics understanding. The two examples I have in mind are magnetic reconnection and the collisionless interaction between a population of energetic particles and a cooler gas or plasma, in particular the interaction between galactic cosmic rays and the interstellar medium.

II. Magnetic Reconnection

Magnetic reconnection is an exceedingly complex phenomenon with many ramifications for astrophysics. I will not attempt to give the theory of reconnection or even to survey all the possible applications, but will simply make a few remarks representing my biases with respect to its treatment.

First, according to Petschek [1] reconnection near a magnetic x point can go only as fast as the Alfvén speed. This is the case even if the layer across which reconnection occurs is so thin that the resistive rate of reconnection is faster. Thus, Petschek's limit on reconnection rates is an upper limit imposed by the inertia of the plasma. Magnetic flux could be destroyed much faster than the Alfvén speed if it were not necessary to move the plasma out of the way.

Second, a great deal of effort has been devoted to establishing that reconnection can actually go as fast as the Alfvén speed. It seems to me that if the condition postulated by Petschek could be achieved, namely a rather narrow x point region outside of which the active field were finite — then Petschek's rate of reconnection is an entirely acceptable estimate. The real trouble seems to hinge on setting up Petschek's conditions. For a simple x point the Alfvén speed near it is, of course, very small so the rate of reconnection is consequently very small.

Thus, the real problem would seem to be to have strong fields near the x point, whose strength can be used to move plasma rapidly through the x point region as Petschek originally conceived the process. This can happen in two ways. One, if there is already a boundary layer as, for example, is the case for the boundary layer between the earth's magnetosphere and the solar wind. Here it seems unquestionable that reconnection can indeed proceed rapidly.

The second way is by a strong MHD instability to occur just prior to reconnection which leads to a very sharp gradient in the magnetic field strength. This would provide a strong magnetic field near the x point to drive rapid reconnection. According to current understanding, this is the two fold mechanism that leads to reasonably rapid reconnection in the internal disruption phenomena in tokamaks. The strong MHD instability is the $m = 1$ kink instability that occurs inside the $q = 1$ surface of the tokamak. In this MHD mode the entire internal region of the tokamak is buckled while the outside region is undisturbed. As a result the magnetic field lines of different topology are pressed together until a strong gradient is established near the $q = 1$ surface, as postulated by Petschek. It is relaxed by the reconnection process as shown by Park [2], and this is the nonlinear process limiting the amplitude of the kink mode.

Thus, the understanding of reconnection processes in many astrophysical instances, in particular the solar flare phenomenon, probably requires the exploration of possible MHD instabilities leading up to the situation where rapid reconnection may occur. This may be as important as further elucidation of the actual details of the reconnection process in the neighborhood of the x point.

Granted the importance of MHD instabilities for reconnection, there is a further problem involved with explosively fast reconnection such as is characterized by solar flares as contrasted to somewhat slower phenomena such as internal disruptions in tokamaks. To understand this point one must introduce into the discussion some generic parameter such as β. As the solar atmosphere and photosphere evolve towards a flare situation, this parameter will change slowly passing through some critical value β_c for which the MHD mode is marginally stable. The evolution of β proceeds at a reasonably slow rate characterized by either resistance or slow motions in the photosphere. The growth rate γ in turn depends on $\beta - \beta_c$, so when β first exceeds β_c, γ will be quite small of order $(d\gamma/d\beta)(\beta - \beta_c)$. The mode will at first grow slowly and eventually reach an amplitude large enough to relax the situation leading to instability in the first place. In other words the MHD mode will fizzle, going off very slowly at a rate only several times the rate $d\beta/dt$, instead of with a bang as is observed. It is probably necessary to have another mode which triggers the MHD mode by suddenly reducing β_c. That is, nonlinear evolution of the two MHD modes is very important for establishing a sudden explosive release of energy. For this case, the reconnection step may be only the last step in the solar flare phenomena.

What MHD instabilities could be involved in solar flares? It seems to me that there is one possible source of instability that may be worth exploring. This is intimately concerned with the sunspot phenomenon itself. In a sunspot an enormous amount of power is bottled up by the suppression of the convective transport of energy by the magnetic fields in the sunspot. This is responsible for the cooling of the sunspot plasma, and finally the lower pressure in the sunspot is responsible for holding the strong magnetic field inside the sunspot. This seems to me to be a situation ripe for instabilities. Further, the amount of power held down by the sunspot fields is many times the energy release in many solar flares. Remember, most solar flares do not emit enough energy to be visible in white light, while most sunspots suppress up to three quarters of the solar radiation that would otherwise pass through these areas. If a fraction of the power confined by sunspot magnetic fields would somehow be released on hydromagnetic time scales, storage of energy would not need to be postulated at all to explain flares!

Although I am not certain, I do not believe such a picture need be in contradiction with observation.

Let me conclude this section by mentioning another situation where rapid reconnection and release of energy could be quite interesting [3]. Let a close pair of binary stars, which are not corotating, each have a magnetic field with some interconnecting lines between them. Then the field in the interstar region would get wound up and tangled at roughly the rotation rate of one star. The field energy will rapidly build up to a value much larger than the vacuum field. This situation should lead to many violent flares and should have very interesting observational consequences.

III. Energetic Particles

It is remarkable in how many different astrophysical situations there are energetic particles. We certainly know that radio emission in every radio source is produced by energetic electrons. We see direct emission or high energy particles by the sun. Perhaps the best known population of high energy particles are the cosmic rays.

These energetic particles form a plasma in their own right, and there are strong interactions between them and the plasma they are embedded in. Over the years people have accepted the existence of these interactions, which are collective plasma interactions primarily through hydromagnetic waves. However, the nonlinear physics of these interactions has not been thoroughly worked out. Because of their importance, I feel it is a place where plasma theory could make a significant contribution to astrophysics.

I will discuss the theory of interactions between the cosmic ray plasma and the interstellar plasma. This is a good illustration of energetic particle plasma interactions.

First we must remember that the interstellar medium is filled with a magnetic field of several microgauss magnitude and more or less uniform on small scales. The gyroradius of a GeV proton in such a field is 10^{12} cm. This is very small compared to the scale of variation of the magnetic field so that, in the absence of interactions, the cosmic ray moves in a helical orbit along the field. It preserves its adiabatic invariant $p_{\perp}^2/B$ and, if the fields are static, its energy.

To determine the type of waves the cosmic ray can interact with, we employ the condition for resonance

$$\omega - k_z v_z = n\Omega \tag{1}$$

where ω and k_z are the wave frequency and projection of the wave number k on $\vec{B}$. Similarily v_z is the projection of the velocity, and $\Omega = cp/eB$ is the relativistic gyrofrequency where p is the total momentum, and n = 0, ±1, ±2 ... . Recall that $v_A = \omega/k_z$ is about 10^7 cm/sec. in the interstellar medium. We note that ω is negligible in Eq. (1) and cosmic rays can interact resonantly with Alfvén waves, satisfying

$$k_z v_z = \Omega \quad . \tag{1a}$$

(For such waves $\omega = k_z v_A = (v_A/v_z)\Omega \approx v_A/c\ \Omega \ll \Omega_o$ if $\gamma \ll c/v_A$ where $\gamma \equiv \varepsilon/mc^2$ is the relativistic factor.)

Because the waves are slow compared to the cosmic rays, the net result of interactions with separate wave packets is a small change in pitch angle of random sign $\Delta\Theta \sim \delta B/B$ with much smaller relative energy change. Therefore the rate of diffusion in pitch angle is

$$\frac{(\Delta\theta)^2}{t} \approx \Omega \frac{(\delta B)^2_{k_z}}{B^2} , \tag{2}$$

where $(\delta B)^2_{k_z}$ is the intensity of fluctuations whose k_z is within a factor two of k_z given by Eq. (1a), Ω/v_z. Because of the large frequency Ω in Eq. (2) the pitch angle scattering of cosmic rays is large even if $(\delta B)^2$ is comparatively small. Thus, one of the primary interactions with background plasma is the isotropization of the cosmic rays.

It is possible that Alfvén waves are already present in the wave length range of Eq. (1a). This is unlikely though because it is difficult to produce waves directly on this scale. Also, because of strong wave damping, the chances that the waves are part of a Kolmogoroff spectrum produced at larger wave length are also small. The most likely source of such waves is generation by the cosmic rays themselves through the same resonance Eq. (1a). If the cosmic rays have any free energy, such as anisotropy in momentum space or a bulk drift through the plasma, it is possible for them to amplify the waves with which they are interacting. In summary, if the cosmic rays have any anisotropy, the free energy associated with this anisotropy goes to generate waves, and the same interaction goes to reduce the anisotropy.

Before pursuing the interaction further, let us consider Eq. (2) a little more carefully near $\theta = \pi/2$. Since from Eq. (1a) $k_z = \Omega/v_z = \Omega/(v \cos \theta)$, k_z approaches infinity as θ approaches $\pi/2$. Therefore it is the shortest waves that scatter a cosmic ray past the $\pi/2$ pitch angle, actually reversing its direction. But as k_z approaches infinity, ω approaches Ω and the waves become strongly damped by ion cyclotron damping of the background plasma. In this same pitch angle scattering range the quasilinear calculation leading to Eq. (2) also breaks down [4]. This is because the change in θ by a single wave packet $\Delta\theta$ may be comparable to or larger than $(\pi/2 - \theta)$. Because the resonance condition is very sensitive to θ near $\pi/2$, such a change is not sufficiently small for a diffusion theory to apply. In other words the cosmic ray must jump over the pitch angle region in a single scattering if the cosmic ray is to reverse direction. This can only be the case if

$$\Delta\theta \approx \left(\frac{\delta B}{B}\right)_{k_z} \gtrsim \frac{\pi}{2} - \theta \tag{3}$$

for some value of θ. If equation (3) is violated for all θ, the cosmic ray cannot reverse its direction. Equation (2) is very crude, but numerical experiments indicate it is approximately correct, and it is very difficult to write down a more precise condition. If Eq. (3) is violated, we can only conclude that the cosmic rays are isotropic separately in the region $0 < \theta < \pi/2$ and $\pi/2 < \theta < \pi$, but we can draw no conclusion as to the relative values of the distribution function in both regions.

We now consider the problem of determining the amplitudes of the waves if they are self-generated by the cosmic rays. Imagine sources and sinks of the cosmic rays along a single line of force. Then because of the interaction with the background plasma, they cannot move freely at the speed of light from source to sink. They will be slowed down to some drift speed v_D, and consequently their density or pressure will vary along the line. Our goal shall be to write down equations relating v_D and L, the scale length for cosmic ray pressure variation. A strong interaction will lead to a small drift velocity and a long time for the cosmic ray to pass from source to sink.

Since v_D corresponds to free energy of the cosmic rays, they will amplify a resonant wave of wave number k with growth rate [5]

$$\Gamma \approx \left(\frac{n_{cr}}{n}\right)_p \left(\frac{v_D - v_A}{v_A}\right) \Omega_o \ , \tag{4}$$

where $(n_{cr})_p$ is the number of cosmic rays in resonance with the wave; i.e., those with p within a factor 2 of eB/kc. Ω_o = eB/mc is the nonrelativistic ion cyclotron frequency. Γ is clearly proportional to the density of cosmic rays and increases linearly with the bulk velocity v_D. If v_D is less than v_A, then the waves are damped.

The waves are amplified by the cosmic rays and damped by normal damping processes. It is important to decide the most important such damping process which in turn depends on the nature of the interstellar medium.

The simplest damping process occurs when the background plasma is partially ionized. In this case charge exchange produces a transfer of momentum between the ions and neutrals. If the frequency of the wave is larger then this transfer rate, the neutrals do not move with the wave, and the wave damps at just half this transfer rate

$$\Gamma^* = \frac{n_o \langle \sigma v \rangle}{2} \approx 10^{-9} \ n_o \ s^{-1} \tag{5}$$

where n_o is the neutral density.

Under these conditions it is easy to see what must happen in our model problem. The flow of cosmic rays from source to sink at first leads to a large v_D. By (4) this leads to amplification of the waves at a rate $\Gamma >> \Gamma^*$. Eventually the waves grow to such an amplitude that pitch angle scattering becomes rapid enough to reduce v_D. v_D then reduces to a small enough value that Γ and Γ^* given by Eq. (5) just balance, and the wave intensity becomes a constant. We can estimate the resulting values of v_D and $\delta B/B$ if we make use of the approximate relation

$$\frac{(\Delta\theta)^2}{\varepsilon}\left(\frac{v_D - v_A}{v_A}\right) \approx \beta\, \Omega\left(\frac{\delta B}{B}\right)^2 (v_D - v_A) \approx \frac{c^2}{L} \tag{6}$$

derivable from elementary transport theory, where β is a number representing the effect of the average over pitch angles. If Eq. (3) is only just satisfied, β may be considerably smaller than unity. Setting $\Gamma = \Gamma^*$, we have with Eq. (6) two equations for the unknowns $(\delta B/B)^2$ and $v_D - v_A$. Solving we have

$$v_D = \left(\frac{n}{n_{cr}}\right) \frac{\Gamma^*}{\Omega_o} v_A + v_A \tag{7}$$

$$\left(\frac{\delta B}{B}\right)^2 = \frac{c^2}{\beta v_A L\Gamma^*} \left(\frac{n_{cr}}{n}\right)\left(\frac{\Omega_o}{\Omega}\right) \quad . \tag{8}$$

Of course, these relations involve k and p parametrically and these are related by the resonance condition. Note that the drift velocity is independent of L and is inversely proportional to n_{cr}. That is to say, the fewer the number of cosmic rays the larger v_D must be to achieve marginal stability. $(\delta B/B)^2$ must adjust to maintain v_D at (7) in the presence of the density gradient n_{cr}/L. The observed value of n_{cr} is proportional to $\gamma^{-1.5}$ so $v_D \sim \gamma^{1.5}$, $(\delta B/B)^2 \sim \gamma^{-.5}$.

Equations (7) and (8) apply to a partially ionized plasma. If the plasma is nearly fully ionized, Γ^* goes zero and Eq. (8) leads to such large amplitudes that nonlinear damping exceeds Γ^* and we must balance Γ against $\Gamma_{n\ell}$, the principal nonlinear damping mechanism. Again there are several damping processes such as mode decay and nonlinear damping. If $v_i > v_A$ where $v_i = (T/M)^{1/2}$ is the thermal velocity, then only nonlinear Landau damping of Alfvén waves exists.

This damping process has been evaluated by a number of people [6]. It is a process where two Alfvén waves interact to produce a nonresonant beat wave with variable magnetic field strength. This variable field in turn acts on ions resonant with the beat wave by means of their magnetic moment. The result is a transfer of energy from the higher frequency Alfvén wave to the lower frequency one with some energy going to the resonant ions. In this way, energy is removed from the Alfvén waves so that the process can be interpreted as nonlinear damping.

The magnitude of this damping is easy to estimate and yields the approximate result

$$\Gamma_{n\ell} = \left(\frac{\pi}{2}\right)^{1/2} \frac{\langle v_\perp^4 \rangle}{16 v_i^4} (k_1 - k_2) v_i \frac{(\delta B_2)^2}{B^2} \tag{9}$$

for damping of wave 1 by wave 2 of lower frequency. δB_2^2 is again the magnetic fluctuation intensity within a factor two of $k_2 < k_1$. We have written the average over $\langle v_\perp^4 \rangle$ in place of its value $8 v_i^4$ for a Maxwell distribution because not all the resonant ions are always able to participate fully in the damping. This is because, if the beat wave fluctuation in magnitude of the magnetic field is too strong, the resonant ions may become trapped in less time than the decorrelation time of the beat wave, and the simple expression of (9) is an overestimate. In fact, this is often the case. [7]

Techniques for correcting expression (9) have not yet developed. However, one may argue roughly that $v_\perp^4$ in the average should be simply multiplied by $\omega_{Dc}/\omega_{trap}$ [8] where

$$\omega_{Dc} \approx k_1(v_\phi - v_g) \approx \frac{kv_\perp^2}{2c} \tag{10}$$

is the rate of decorrelation [9], and

$$\omega_{trap} = \left(k \frac{v_\perp^2}{B} \delta|B|\right)^{1/2} \approx kv_\perp \frac{\delta B_1}{B} \tag{11}$$

is the rate of trapping. We use $\delta|B| = \delta B_1\ \delta B_2/B$ and treat δB_1 and δB_2 as comparable for this very simplified estimate.

Simply for the purposes of exhibition we consider two limiting cases.

In the first case (I) let $\omega_{trap} < \omega_{Dc}$ be satisfied for all $v_\perp$ so that the normal nonlinear Landau damping theory applies. Replacing $\delta|B| \approx \delta B\ \delta B_2/B$ $\langle v_\perp^4\rangle$ by 8 v_i^4 , we have

$$\Gamma^I_{n\ell} \approx \left(\frac{\pi}{8}\right)^{1/2} k\ v_i \left(\frac{\delta B}{B}\right)^2 \tag{12}$$

where we set $(\delta B)_2^2 \approx (\delta B)^2$ and $k_1 - k_2 \approx k$. [By (10) and (11) case (I) applies when $\delta B/B < v_i/c$, and is the weak fluctuation limit.] Setting Γ from (4) equal to $\Gamma^I_{n\ell}$ and employing (1a) and (6) we can again solve for v_D and $\delta B/B$. There results

$$v_D^I = \left(\frac{\pi}{8}\right)^{1/4}\left(\frac{n}{n_{cr}}\right)^{1/2}\left(\frac{v_A v_i c}{\beta I \Omega_o}\right)^{1/2} + v_A \tag{13}$$

$$(\frac{\delta B}{B})^2_I = (\frac{8}{\pi})^{1/4}(\frac{n_{cr}}{n})^{1/2}(\frac{c^3\Omega_o}{v_A v_i \Omega^2 \beta L})^{1/2} \quad . \tag{14}$$

Note that v_D now depends on the scale length L. It must be realized that v_D is a function of the energy parameter γ as well as position. It is determined by (13) once n_{cr} as a function of γ is known. This in turn is determined by (13) together with the specifications of sources and sinks. If we take the observed value of $n_{cr} \sim \gamma^{-1.5}$, we find for $v_D >> v_A$, $v_D \sim \gamma^{.75}$ increases with energy. Similarily $(\delta B/B)^2$ is a function of k. This dependence is found from the dependence of (14) on γ and using Eq. (1a). Hence if $n_{cr} \sim \gamma^{-1.5}$, $(\delta B/B) \sim \gamma^{1/4} \sim k^{-1/4}$. Waves interacting with lower energy cosmic rays are the ones that satisfy the condition of case (I).

Case (II) is the opposite limit to that of case (I). For case (II) $\omega_{trap} > \omega_{Dc}$ for all $v_\perp$. Then we have

$$\Gamma^{II}_{nL} = (\frac{\pi}{8})^{1/2} \frac{\langle v_\perp^3 \rangle}{16 v_i^2 c} (\frac{\delta B}{B}) \; kv_i = \frac{3\pi}{64} \frac{kv_i^2}{c} (\frac{\delta B}{B}) \quad . \tag{15}$$

Again setting $\Gamma = \Gamma^{II}_{n\ell}$ and employing (1a) and (6), we have

$$v_D^{II} = (\frac{3\pi}{64})^{2/3} (\frac{v_A v_i^2 \Omega^{1/2}}{c\beta^{1/4} L^{1/2} \Omega_o})^{2/3} (\frac{n}{n_{cr}})^{2/3} + v_A \; , \tag{16}$$

$$(\frac{\delta B}{B})^2_{II} = (\frac{64}{3\pi})^{2/3} (\frac{c^4 \Omega_o}{v_i^2 v_A \Omega^2 L\beta})^{2/3} (\frac{n_{cr}}{n})^{2/3} \quad . \tag{17}$$

If we take $n_{cr} \sim \gamma^{-1.5}$, and $v_D >> v_A$, $v_D \sim \gamma^{2/3}$ and $(\delta B/B)^2 \sim \gamma^{+1/3} \sim k^{-1/3}$ for case II.

Equations (13) and (16) would solve our problem of determining the relation between the scale size for pressure variation of the cosmic rays L and their bulk drift velocity if the assumptions could be justified and the constants such as β better determined. The question of which limiting case applies depends on which inequality in $\omega_{Dc} \gtrless \omega_{trap}$ applies, the upper inequality applying to case (I). Using Eq. (10) through (11) this reduces to

$$\left(\frac{\delta B}{B}\right)^2 \lessgtr \frac{v_i^2}{4c^2} , \qquad (18)$$

and $(\delta B/B)^2$ is determined by (14) or (17) accordingly. However, in case (I) $\Gamma^{I}_{n\ell} < \Gamma^{II}_{n\ell}$ since $\omega_{trap} << \omega_{Dc}$, while in case (II) the opposite inequality prevails. Thus, in either case the smaller result of the (13) and (16) automatically applies. Similarily the larger of (14) and (17) applies.

In addition to the physical uncertainties involving the pitch angle scattering through $\pi/2$ and the saturation of the damping, I have also simplified the problem by neglecting the coupling between waves with different k's and cosmic rays with different p's. Thus, I have solved algebraic equations rather than integral equations which lead to a more correct theory. There seems little point into going into these complications until the basic uncertainties are removed.

I have emphasized the control of the background plasma over the cosmic rays by calculating v_D assuming the background plasma parameters as known, and by even setting its velocity to zero. Often the pressure of the energetic component can be significant and it can play a dynamic role in the motion of the plasma. The pressure gradient will exert a force on the waves which in turn through their damping will transfer motion to the plasma. At the same time the cosmic rays can exchange energy with the plasma, which can be of considerable importance for applications.

IV. Conclusion

In conclusion I would like to mention some phenomena to which the physics of Sec. III should be applicable.

One of the most important of these is the propagation of cosmic rays through the galaxy. It is accepted that cosmic rays are galactic in origin, and it is known from studies of their composition that they leave the galactic disk within at most a few million years after their generation. The details of their motion cannot be known until the galactic magnetic field is better understood. However if the magnetic field is open to the intergalactic medium, then the picture sketched in Sec. III of a cosmic ray pressure gradient over some length L should be relevant. Then making use of known properties of the interstellar medium, measured values for the cosmic ray anisotropy and spectrum, and the dependence of cosmic ray age on energy, it should be possible to derive important properties concerning the sources of generation of cosmic rays.

A second application of these ideas is to the proposed origin of cosmic rays in supernova explosions [10]. If the cosmic rays are produced in the explosion or in the remnant star, again a density gradient will be set up and cosmic rays will drift relative to the expanding supernova envelope with velocity v_D. Now if v_D is small compared to the rate of expansion of the envelope, the cosmic rays are not able to outrun the envelope, and they will be strongly adiabatically decompressed losing a large portion of their energy. Thus, it is important to determine v_D in order to test out this origin theory for cosmic rays.

A third application is to acceleration of cosmic rays by a strong shock [11]. Again the interaction between the shocked fluid and the cosmic rays can force the cosmic rays to cross the shock many times, suffering first order

Fermi acceleration on each pass through the shock. To be sure of this process, it is necessary to work out the details of this interaction along lines suggested in Sec. III.

As a final example, the large radio sources seem to derive their energy from small central sources of energetic particles. These particles would be prevented from flowing out from these sources at high velocities. They would be restricted again to a velocity v_D relative to the intergalactic medium. If v_D is small, it might be the case that the intergalactic medium itself may be accelerated to high velocities outward from the central source.

These examples should make clear the urgency of using the techniques of plasma physics to solve the different problems associated with energetic particle-background plasma interactions.

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References

1. Petschek, H. E., Proceedings of the AAS-NASA Conference on Plasma Physics of Solar Flares NASA sp 50, 425 (Goddard 1963).
2. Park, W., Monticello, D. A., and White, R. B., Nucl. Fusion 20, 1181 (1980).
3. Bahcall, J. B., Kulsrud, R. M., and Rosenbluth, M. N., Nature Physical Sciences 243, 27 (1973).
4. Kaiser, T. B., Jones, F. C., and Birmingham, T. J., Astrophys. J. 180, 239 (1973).
5. Kulsrud, R. M., and Pearce, W. D., Astrophys. J. 156, 445 (1969).
6. Kulsrud, R. M., Important Advances in XXth Century Astronomy" Copenhagen University Observatory (1978) p. 317.
7. Cesarsky, C. J., and Kulsrud, R. M., "Origin of Cosmic Rays" G. Setti G. Spada AW Wolfendale (Eds) 251 (1981).
8. Kulsrud, R. M., Proc. of the 16th Int'l. Conf. on Cosmic Rays (Tokyo) 12, 128 (1979).
9. Foote, E. A., and Kulsrud, R. M., Astrophys. J. 233 302 (1979).
10. Kulsrud, R. M., and Zweibel, E. G., Proc. of the 14th Int'l. Conf. on Cosmic Rays (Munich) 1975 (OG9-1-12).
11. Blandford, R. D., and Ostriker, J. P., Astrophys. J. 221, 37 (1978).