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Generation of Mesoscale Magnetic Fields and the Dynamics of Cosmic Ray Acceleration

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Abstract

The problem of the cosmic ray origin is discussed in connection with their acceleration in supernova remnant shocks. The diffusive shock acceleration mechanism is reviewed and its potential to accelerate particles to the maximum energy of (presumably) galactic cosmic rays (10^{18} eV) is considered. It is argued that to reach such energies, a strong magnetic field at scales larger than the particle gyroradius must be created as a result of the acceleration process, itself. One specific mechanism suggested here is based on the generation of Alfven wave at the gyroradius scale with a subsequent transfer to longer scales via interaction with strong acoustic turbulence in the shock precursor. The acoustic turbulence in turn, may be generated by Drury instability or by parametric instability of the Alfven waves.

Keywords:

acceleration of particle, cosmic ray, magnetic field, shock wave, MHD, turbulence

1. Introduction

Cosmic rays (CRs) were discovered almost a century ago, yet their origin is unknown. There is sufficient evidence that at least part of their spectrum ($E < 10^{18}$ eV) originates in the Galaxy, while particles of higher energies are thought to come to us from extragalactic sources. One simple theoretical argument is that particles with such high energies (usually referred to as ultra high energy CRs, UHECR) could not be confined to the Galaxy, on account of their large gyroradius. Second, their (power law- $E^{-2.7}$) spectrum is harder than that of (presumably) galactic particles at $E < 10^{18}$ eV ($E^{-3.1}$), which is consistent with their extra-galactic origin. The last point becomes clear if we turn to the other break on the overall CR spectrum, the one at $E \sim 10^{15}$ eV, commonly known as the "knee", shown in Fig. 1. The spectrum above this energy steepens (from $E^{-2.7}$ to $E^{-3.1}$), so that the premise of their extragalactic origin would require an explanation of why the galactic part of the spectrum terminates, while the extra-galactic part appears *exactly* at the knee energy. As we shall see, to explain the cosmic ray power law spectrum between the "knee" at 10^{15} eV and the "ankle" at 10^{18} eV in terms of acceleration within the Galaxy is one of the most serious challenges of contemporary acceleration mechanisms and one of the main motivations of this study.

The explanation of the spectrum beyond the ankle (the highest energy event observed so far is $3 \cdot 10^{20}$ eV) poses a major challenge to fundamental physics. Given the distance of possible accelerators (at least a few tens of Mpc), particles of such high energy should have experienced significant losses through their interaction with the cosmic microwave background radiation (the so called Greisen, Zatsepin, Kusmin or GZK cut-off) while propagating over long distances. Note

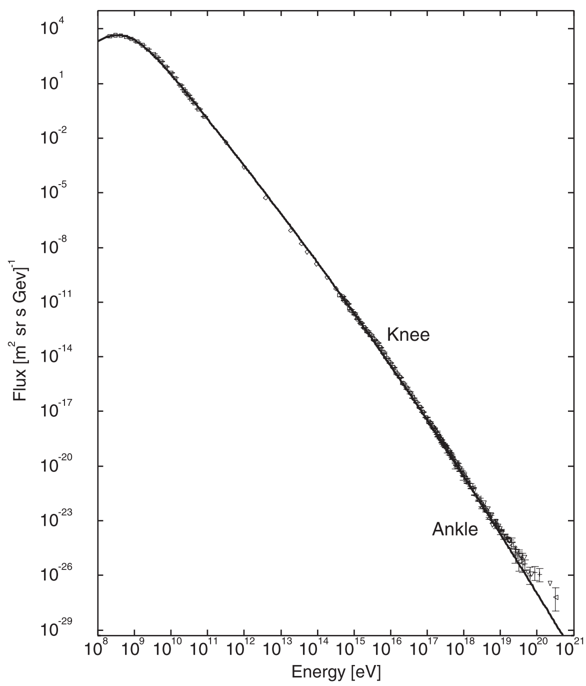

Fig. 1 Cosmic Ray spectrum.

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that UHECR are sub-atomic particles with the energy of a well-hit baseball. This energy exceeds (by at least three orders of magnitude) that achievable by all existing terrestrial accelerators (e.g., Large Hadron Collider, LHC-10^{17} eV). Another fundamental aspect of the problem of CR origin is that they are the ultimate receptacle of a significant portion of the gravitational energy in the Universe. Indeed, star formation from the gravitational collapse of the primordial gas with subsequent SN (supernova) explosions and their blast waves results in CR acceleration. On the observational side, there is an evidence [1-3] (in the form of both synchrotron and inverse Compton radiation) that electrons of energies up to 100 TeV are accelerated in the supernova shock waves.

2. Acceleration mechanism

The leading CR acceleration mechanism, namely the diffusive shock acceleration (DSA, also known as the 1st -order Fermi mechanism) was proposed originally by Fermi in Ref. [4], and in its modern form by a number of authors in the late seventies [5-8]. The mechanism is basically simple – particles gain energy by bouncing between converging upstream and downstream regions of the flow near a shock wave such as that from an SNR (supernova remnant) shock. This mechanism requires magnetic fields. First, the field binds particles to the accelerator (shock wave), in the direction perpendicular to the field. Confinement in the direction *along* the field lines is, in turn, ensured by particles themselves through the generation of Alfven waves by accelerated particles streaming ahead of the shock. This occurs via Doppler resonance, $\Omega = k(p/m)\mu$, where Ω and p are the (nonrelativistic) gyrofrequency and momentum, k is the wave number, m is the particle mass, and μ is the cosine of its pitch angle. These waves, in turn, scatter particles in pitch angle (at the rate, $\nu \sim \Omega(mc/p)\ (\delta B/B_0)^2$, where c is the speed of light) back and force so that they can gain energy by repeatedly crossing the shock. Particle self-confinement along the field is thus diffusive and the diffusivity $\kappa \sim c^2/\nu$ is inversely proportional to the fluctuation energy $\boldsymbol{\delta B}^2$, as the fluctuating field is responsible for pitch-angle scattering. However, the mean field B_0 ultimately determines the acceleration rate and the particle maximum energy since it sets the work done by the induced electric field on the particles. Indeed, $E_{max} \sim (e/c)u_s B_0 R_s$, where u_s and R_s are the speed and typical size of the shock wave, such as the radius of the SNR shock. The fluctuating part, δB, is typically assumed to be saturated at most at the level $\delta B \sim B_0$, which provides pitch-angle scattering at the rate of gyrofrequency, and thus limits the particle mean free path (m.f.p.) along the field to a distance of the order of gyroradius (the so called Bohm diffusion limit). An important thing to keep in mind is that, due to the resonance condition $kp = const$, confinement of particles of higher energies requires that longer waves need be excited.

The most critical test of this mechanism is the requirement that it accelerates galactic CRs to the energy of 10^{15} eV over the life-time of supernova remnant shocks. Even with the above "optimistic" estimates of the turbulence level, the mechanism passes this test at best only marginally. If the turbulence level is lower, then the maximum energy should be reduced proportionally. There are indeed a number of phenomena which may either reduce the turbulence level [9], or which can shift the turbulence spectrum (in wave number) away from resonance with the high energy particles and therefore cause their losses [10].

Another reason for concern about this mechanism, at least in its standard (Bohm limit) version, is its seeming inability to explain acceleration of particles beyond 10^{15} eV. As was discussed above, the cosmic ray spectrum has only a break at this energy, and continues to about 10^{18} eV where the extragalactic component is believed to start dominating the spectrum.

One approach to this problem is to generate a fluctuating component δB significantly exceeding the unperturbed field B_0 [11]. Physically, such generation is possible since the free energy source is the pressure gradient of accelerated particles, which in turn may reach a significant fraction of the shock ram energy. Specifically, the wave energy density $(\delta B/B_0)^2$ may be related to the partial pressure P_c of CRs that resonantly drives these waves through the relation [12]

$$\left(\delta B / B_0\right)^2 \sim M_A P_c / \rho u_s^2 \qquad (1)$$

where $M_A = u_s/V_A \gg 1$is the Alfven Mach number and ρu_s^2 is the shock ram pressure. Of course, when $\delta B/B_0$ exceeds unity, particle dynamics, and thus their confinement and acceleration rates, are very difficult to assess if the turbulence spectrum is sufficiently broad. The numerical studies by [11] showed that at least in the case of an MHD (magnetohydrodynamic) description of the background plasma and rather narrow wave (and particle energy) band, the amplitude of the principal mode can reach a few times that of the background field. Moreover, the authors of Ref. [13] argue that in the case of efficient acceleration, field amplification may be even stronger, reaching a mG (10^{-3} Gauss) level from the background of a few μG (10^{-6} Gauss) ISM field, thus providing acceleration of protons up to 10^{17} eV in SNRs.

Recently, the authors of Ref. [14] approached this problem from a different perspective. They considered a Kolmogorov turbulent cascade to small scales assuming the waves are generated by efficiently accelerated particles on the long-wave part of the spectrum. They obtained a particle maximum energy similar to that of [13].

Apart from the excitation of magnetic fluctuations during acceleration process, there is yet another aspect of the CR-magnetic field connection discussed in the literature. Zweibel [15] points out that since CRs were already present in young galaxies (observed at high redshifts), magnetic fields of appreciable strength must also have been there at that time. She emphasizes, however, that the approximate equipartition between the CR and magnetic field energy established by the current epoch in our Galaxy is not required for the acceleration mechanism and presumably results from the fact that they both have a common energy source, namely the supernovae.

Indeed, as we discussed already, the magnetic field strength merely determines the *maximum* energy of accelerated particles, given the time available for acceleration and the size of the accelerator. If the latter are sufficient then the total energy of accelerated particles can, and in most of the DSA models *does*, exceed that of the magnetic field. The latter, in turn, remains unchanged, apart from the conventional compression at the shock and the MHD fluctuations discussed earlier.

In this paper we discuss the possibility of a *different* scenario, in which the magnetic field may absorb a significant part of the shock energy as a result of the acceleration process, which may in fact be strongly enhanced. The mechanism of such enhancement is based on the transfer of magnetic energy to longer scales, which we call *inverse cascade* for short, even though specific mechanisms of such transfer may differ from what is usually understood as a cascade in MHD turbulence. This transfer is limited only by some outer scale L_{out} such as the shock precursor size $\kappa(p_{max})/u_s \sim r_g(p_{max})\,c/u_s \gg r_g(p_{max})$. This approach is in contrast to the above discussed models [13,14], which operate on the generated magnetic fields with the scale lengths of the order of the Larmor radius $r_g(p_{max})$ of the highest energy particles and smaller. The advantage of the inverse cascade for the acceleration is that the turbulent field at the outer scale $\delta B(L_{out}) \equiv B_{rms}$ (which necessarily must have long autocorrelation time) can be obviously regarded as an "ambient field" for accelerated particles of all energies. If $B_{rms} \gg B_0$, then the acceleration can be enhanced by a factor B_{rms}/B_0. Note that the resonance field $\delta B(r_g)$ may remain smaller than B_{rms}, so that standard arguments about Bohm diffusion apply, and it is less likely that this field will be rapidly dissipated by nonlinear processes, such as induced scattering on thermal protons, not included in the enhanced acceleration model [11,13].

As it should be clear from the above, an adequate description of the acceleration mechanism must include *both* particle and wave dynamics on an equal footing. In fact the situation is even more difficult, since the acceleration process turns out to be so efficient that the pressure of accelerated particles markedly modifies the structure of the shock (both the overall shock compression and the flow profile).

3. Accelerated particles and plasma flow near the shock front

The transport and acceleration of high energy particles (CRs) near a CR modified shock is usually described by the diffusion-convection equation. It is convenient to use a distribution function $f(p)$ normalized to $p^2 \mathrm{d}p$.

$$\frac{\partial f}{\partial t} + U\frac{\partial f}{\partial x} - \frac{\partial}{\partial x}\kappa\frac{\partial f}{\partial x} = \frac{1}{3}\frac{\partial U}{\partial x}\,p\frac{\partial f}{\partial p} \qquad (2)$$

Here x is directed along the shock normal, which for simplicity, is assumed to be the direction of the ambient magnetic field. The two quantities that control the acceleration process are the flow profile $U(x)$ and the particle diffusivity $\kappa(x, p)$. The first one is coupled to the particle distribution f through the equations of mass and momentum conservation

$$\frac{\partial}{\partial t}\rho + \frac{\partial}{\partial x}\rho U = 0 \qquad (3)$$

$$\frac{\partial}{\partial t}\rho U + \frac{\partial}{\partial x}\left(\rho U^2 + P_c + P_g\right) = 0 \qquad (4)$$

where

$$P_c(x) = \frac{4\pi}{3}mc^2\int_{p_{inj}}^{\infty}\frac{p^4\mathrm{d}p}{\sqrt{p^2+1}}f(p,x) \qquad (5)$$

is the pressure of the CR gas and P_g is the thermal gas pressure. The lower boundary p_{inj} in momentum space separates CRs from the thermal plasma that is not directly involved in this formalism but rather enters the equations through the magnitude of g at $p = p_{inj}$, which specifies the injection rate of thermal plasma into the acceleration process. The particle momentum p is normalized to mc. The spatial diffusivity κ, induced by pitch angle scattering, prevents particle streaming away from the shock, thus facilitating acceleration by ensuring the particle completes several shock crossings.

The system (2-5) indicate a marked departure from the test particle theory. Perhaps the most striking result of the nonlinear treatment is the bifurcation of shock structure (in particular shock compression ratio) in the parameter space formed by the injection rate, shock Mach number and particle maximum momentum [16].

4. Wave dynamics in the CR shock precursor

The transformation of magnetic energy to longer scales, while bearing certain characteristics of the conventional turbulent dynamo problem, is still rather different from it, in its conventional form. First, this process should take place in the strongly compressible fluid near the shock. Second, the Alfven wave turbulence is generated by accelerated particles via Cerenkov emission, and thus is strongly coupled to them. Third, the shock precursor itself is unstable to emission of acoustic waves. The latter phenomenon is known as the Drury instability and will be discussed later. Acoustic waves, in turn, interact with particle generated Alfvenic turbulence, stimulating decay instability (i.e., "inverse cascade").

The spatial structure of an efficiently accelerating shock, i.e., the shock that transforms a significant part of its energy into accelerated particles, is very different from that of the ordinary shock, Fig. 2. Most of the shock structure consists of a precursor formed by accelerated CRs diffusing ahead of the shock. If the CR diffusivitity $\kappa(p)$ depends linearly on particle momentum p (as in the Bohm diffusion case), then, at least well inside the precursor, the velocity profile $U(x)$ is approximately a *linear* function of x, where x points along the shock normal [17]. Ahead of the shock precursor, the flow

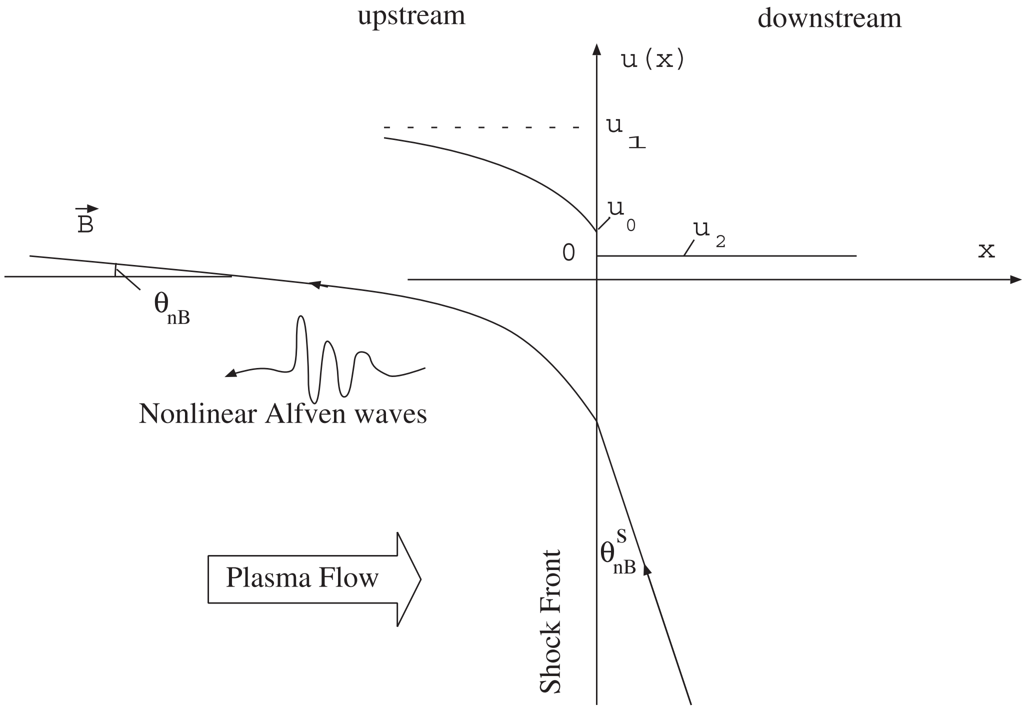

Fig. 2 Schematic representation of nonlinearly accelerating shock. Flow profile with a gradual deceleration upstream is also shown at the top.

velocity tends to its upstream value, U_1, while on the downstream side it undergoes a conventional plasma shock transition to its downstream value U_2 (all velocities are taken in the shock frame). This extended CR precursor (of the size $L_{CR} \sim \kappa(p_{max})/U_1$) is the place where we expect turbulence is generated by the CR streaming instability and where it cascades to longer wavelengths.

4.1 Alfvenic turbulence

The growth rate of the ion-cyclotron instability is positive for the Alfven waves traveling in the CR streaming direction i.e., upstream, and it is negative for oppositely propagating waves. The wave kinetic equation for both types of waves can be written in the form

$$\frac{\partial N_k^{\pm}}{\partial t} + \frac{\partial \omega^{\pm}}{\partial k}\frac{\partial N_k^{\pm}}{\partial x} - \frac{\partial \omega^{\pm}}{\partial x}\frac{\partial N_k^{\pm}}{\partial k} = \gamma_k^{\pm} N_k^{\pm} + C^{\pm}\left\{N_k^{+}, N_k^{-}\right\} \qquad (6)$$

Here $N_k^{\pm}$ denotes the number of quanta propagating in the upstream and downstream directions, respectively. Also, $\omega^{\pm}$ are their frequencies, $\omega^{\pm} = kU \pm kV_A$, where V_A is the Alfven velocity. The linear growth rates $\gamma^{\pm}$ are nonzero only in the resonant part of the spectrum, $kr_g(p_{max}) > 1$. In the most general case, the last term on the r.h.s. represents nonlinear interaction of different types of quanta N_k^+ and N_k^- and, if compressibility effects are present, also interactions between the same type. As seen from this equation, the coefficients in the wave transport part of this equation (l.h.s.) depend on the parameters of the medium through U and V_A, which in turn, may be subjected to perturbations. This usually results in parametric phenomena [18]. We will concentrate on the acoustic type perturbations (which may be induced by Drury instability), so that we can write for the density ρ and velocity U

$$\rho = \rho_0 + \tilde{\rho}; U = U_0 + \tilde{U}$$

The variation of the Alfven velocity $\tilde{V}_A = V_A - V_{A0}$

$$\tilde{V}_A \simeq -\frac{1}{2} V_A \frac{\tilde{\rho}}{\rho_0}.$$

For simplicity, we assume that the plasma $\beta < 1$ (which is not universally true in the shock environment) and neglect the variation of U compared to that of V_A in eq. (6). The above perturbations of V_A in turn induce perturbations of $N_k^{\pm}$, so we can write

$$N_k^{\pm} = \left\langle N_k^{\pm} \right\rangle + \tilde{N}_k^{\pm}$$

Our goal is to obtain an evolution equation for the averaged number of plasmons $\langle N_k^{\pm} \rangle$. Averaging eq. (7) we have

$$\frac{\partial}{\partial t}\left\langle N_k^{\pm}\right\rangle + \left(U \pm V_A\right)\frac{\partial}{\partial x}\left\langle N_k^{\pm}\right\rangle - kU_x \frac{\partial}{\partial k}\left\langle N_k^{\pm}\right\rangle + \frac{\partial}{\partial k}\left\langle kV_A \frac{\tilde{\rho}_x}{\rho_0}\tilde{N}_k^{\pm}\right\rangle = \gamma_k^{\pm}\left\langle N_k^{\pm}\right\rangle + \left\langle C\left(N_k^{\pm}\right)\right\rangle \qquad (7)$$

Here the index x stands for the x-derivatives. To calculate the correlator $\left\langle \frac{\tilde{\rho}_x}{\rho_0}\tilde{N}_k^{\pm}\right\rangle$ in the last equation, we expand the r.h.s. of eq. (6) retaining only the main linear part in $\tilde{N}$

$$\gamma_k^{\pm} N_k^{\pm} + C^{\pm}\left\{N_k^{+}, N_k^{-}\right\} \simeq -\Delta\omega_k^{\pm}\, \tilde{N}_k^{\pm} \qquad (8)$$

The time scale separation between the l.h.s. and r.h.s. of eq. (6) suggests that to lowest order, the linear growth γ^+ rate is approximately balanced by the local nonlinear term C^+. Likewise, the linear damping of the backward waves γ^- may be balanced by their nonlinear growth and conversion of the forward waves C^-. Generally, the $\Delta\omega_k^{\pm}$ in eq. (8) is a 2 × 2 matrix operator. If the wave collision term is quadratic in N, then $\Delta\omega_k^{\pm} \cong \gamma_k^{\pm}$. This balance can be established only for the resonant waves ($\gamma^{\pm} \neq 0$), whereas our primary focus will be on the extended longwave interval $k < 1/r_g(p_{max})$ for which $\gamma \approx 0$. In this domain, cascading from the generation region $k > 1/r_g(p_{max})$ takes place and the refraction (last) term on the l.h.s. of eq. (7) plays a dominant role along, with the nonlinear term on the r.h.s.

To calculate the refraction term we write eq. (6), linearized with respect to $\tilde{N}_k^{\pm}$, as:

$$L^{\pm}\, \tilde{N}_k^{\pm} = -kV_A \frac{\tilde{\rho}_x}{2\rho_0}\frac{\partial}{\partial k}\left\langle N_k^{\pm}\right\rangle \qquad (9)$$

where

$$L^{\pm} = \frac{\partial}{\partial t} + \left(U \pm V_A\right)\frac{\partial}{\partial x} - kU_x \frac{\partial}{\partial k} + \Delta\omega_k^{\pm}$$

Solving eq. (9) for $\tilde{N}_k^{\pm}$, from eq. (7) we thus have the following equation for $\langle \tilde{N}_k^{\pm}\rangle$

$$\frac{\partial}{\partial t}\left\langle N_k^{\pm}\right\rangle + U\frac{\partial}{\partial x}\left\langle N_k^{\pm}\right\rangle - kU_x\frac{\partial}{\partial k}\left\langle N_k^{\pm}\right\rangle - \frac{\partial}{\partial k}D\frac{\partial}{\partial k}\left\langle N_k^{\pm}\right\rangle = \gamma_k^{\pm}\left\langle N_k^{\pm}\right\rangle + \left\langle C\left(N_k^{\pm}\right)\right\rangle \qquad (10)$$

Here we introduced a diffusion operator for the Alfven waves in k space due to *random* refraction by the acoustic perturbations $\tilde{\rho}$ (via the density dependence of V_A), i.e.,

$$D_k = \frac{1}{4}k^2 V_A^2 \left\langle \frac{\tilde{\rho}_x}{\rho_0} L^{-1} \frac{\tilde{\rho}_x}{\rho_0}\right\rangle \qquad (11)$$

D_k is an example of the well-known phenomenon of induced diffusion. Transforming to Fourier space, we first represent $\tilde{\rho}$ as

$$\tilde{\rho} = \sum_q \rho_q e^{\mathrm{i}qx - \mathrm{i}\Omega_q t}$$

and note that due to the local Galilean invariance of L, we can calculate its Fourier representation in the reference frame moving with the plasma at the speed $U(x)$ as:

$$L_{k,q}^{\pm} = \pm \mathrm{i}qV_A + \Delta\omega_k^{\pm} - kU_x \frac{\partial}{\partial k} \qquad (12)$$

Then, eq. (11) can be re-written as:

$$D_k = \frac{1}{2}k^2 V_A^2 \sum_q q^2 \left|\frac{\rho_q}{\rho_0}\right|^2 \Re L_{k,q}^{-1} \qquad (13)$$

The last (wave refraction) term on the r.h.s. of eq. (12) can be estimated as $U_1^2/\kappa(p_{max})$, which is the inverse acceleration time and can be neglected as compared to the frequencies qV_A and $\Delta\omega$. Hence, $\Re L_{k,q}^{\pm-1}$ for we have:

$$\Re L_{k,q}^{\pm -1} \approx \frac{\Delta\omega_k^{\pm}}{q^2 V_A^2 + \Delta\omega_k^{\pm 2}}$$

For further convenience, we introduce here the number of phonons

$$N_q^s = \frac{W_q}{\omega_q^s}$$

where W_q is the energy density of acoustic waves (with $\omega_q^s = qC_s$ as their frequency).

$$W_q = C_s^2 \frac{\rho_q^2}{\rho_0}$$

For D_k in eq. (10) we thus finally have

$$D_k = \frac{k^2 V_A^2}{2C_s^2 \rho_0}\sum_q q^2 \omega_q^s \frac{\Delta\omega_k^{\pm}}{q^2 V_A^2 + \Delta\omega_k^{\pm 2}} N_q^s$$

Note that D_k represents the rate at which the wave vector of the Alfven wave random walks due to stochastic refraction. Of course, such a random walk necessarily must generate larger scales (smaller k), thus in turn facilitating the confinement (to the shock) of higher energy particles. Thus, confinement of higher energy particles is a natural consequence of Alfven wave refraction in acoustic wave generated density perturbations.

4.2 Acoustic turbulence

Unlike the Alfvenic turbulence that originates in the shock precursor from accelerated particles, there are two separate sources of acoustic perturbations. One is related to parametric [18] processes undergone by the Alfven waves in the usual form of a decay of an Alfven wave into another Alfven wave and an acoustic wave. The other source is the pressure gradient of CRs, which directly drives instability. The latter leads to emission of sound waves due to the Drury instability. By analogy with eq. (7) we can write the following wave kinetic equation for the acoustic waves:

$$\frac{\partial}{\partial t}N_q + U\frac{\partial}{\partial x}N_q - qU_x\frac{\partial}{\partial k}N_q = \left(\gamma_q^d + \gamma_D\right)N_q + C\left\{N_q\right\}$$

Here γ_D is the Drury instability growth rate and γ_q^d is that of the decay instability. We first consider the decay instability of Alfven waves. Note, however, that the combination of Drury instability and decay instability can lead to generation of mesoscale fields at a faster than – exponential rate, by coupling together the Drury and decay instability processes.

4.2.1 Decay instability

The mechanism of this instability is the growth of the density (acoustic) perturbations due to the action of the ponderomotive force from the Alfven waves. This force can be regarded as a radiative pressure term appearing in the hydrodynamic equation of motion for the sound waves (written below in the comoving plasma frame)

$$\frac{\partial V}{\partial t} = -\frac{1}{\rho_0}\frac{\partial}{\partial x}\left(C_s^2 \tilde{\rho} + P_{rad}\right)$$

Eliminating velocity by making use of continuity equation

$$\frac{\partial \tilde{\rho}}{\partial t} + \rho_0 \frac{\partial V}{\partial x} = 0,$$

we obtain

$$\frac{\partial^2 \tilde{\rho}}{\partial t^2} = \frac{\partial^2}{\partial x^2}\left(C_s^2 \tilde{\rho} + P_{rad}\right) \tag{14}$$

The Alfven wave pressure can be expressed through their energy

$$P_{rad} = \sum_k \omega_k \left(\tilde{N}_k^+ + \tilde{N}_k^-\right)$$

Using the relation (9) between the density perturbations and the Alfven waves and separating forward and backward propagating sound waves $\rho^{\pm}$, we can obtain from eq. (14) the following dispersion relation for the acoustic branch

$$\omega^2 - q^2 C_s^2 = q^2 \sum_k \frac{\omega_k}{2\rho_0} \mathrm{i}qkV_A L_{k,q}^{\pm -1} \frac{\partial}{\partial k}\left\langle N_k^{\pm}\right\rangle$$

or on writing $\omega = \pm\, qC_s + \mathrm{i}\gamma^{\pm}$, we have the following growth rate of acoustic perturbations

$$\gamma^{\pm} = \frac{q^2}{4\rho_0}\frac{V_A}{C_s}\sum_k k\omega_k L_{k,q}^{\pm -1}\frac{\partial}{\partial k}\left\langle N_k^{\pm}\right\rangle$$

Note that the instability requires an inverted population of Alfven quanta. As they are generated by high energy resonant particles in a finite domain of k space, such an inversion clearly can occur.

4.2.2 Drury instability

This instability also leads to efficient generation of acoustic waves and it is driven by the pressure gradient of the CRs in the shock precursor. The growth rate has been calculated in Ref. [19] (see also [20-22]), and can be written in the form:

$$\gamma_D^{\pm} = -\frac{\gamma_C P_C}{\rho\kappa} \pm \frac{P_{Cx}}{C_s \rho}\left(1 + \frac{\partial \ln\kappa}{\partial \ln\rho}\right) \tag{15}$$

Here P_C and P_{Cx} are the CR pressure and its derivative, respectively, and γ_C is their adiabatic index. For an efficiently accelerating shock $\gamma_C \approx 4/3$. Note that we have omitted a term $-U_x$ which is related to a simple compression of wave number density by the flow and should be generally incorporated into the r.h.s. of eq. (14) (see [19]). This term is smaller by a factor C_s/U than the second (destabilizing) term. The first term is damping caused by CR diffusion, calculated earlier by Ptuskin [23].

5. Mechanisms of transfer of magnetic energy to larger scales

As it follows from the above considerations, there are a variety of nonlinear processes that can lead to the transfer of magnetic energy (generated by accelerated particles in form of the resonant Alfven waves) to longer scales. First, as it can be seen from eq. (10) (last term on the l.h.s.) scattering of the Alfven waves in k-space due to acoustic perturbations transfers magnetic fluctuations away from the resonant excitation region to smaller (but also to larger) k. Second, the nonlinear interaction of Alfven waves and magnetosonic waves represented by the wave collision term on the r.h.s. can be responsible for such process. It is also well known that in the presence of nonzero magnetic helicity there is a strong inverse cascade of magnetic energy [24,25]. Finally, even in the frame work of the weak turbulence, induced scattering of Alfven waves on thermal protons leads to a systematic decrease in the energy of quanta which, given the dispersion law, means again transformation to longer waves.

Returning to the wave refraction process on acoustic perturbations generated by the Drury instability, it is important to emphasize the following. As is seen from the instability growth rate, eq. (15), it is proportional to the gradient of P_c. As the latter should be increased as a result of instability, through a better confinement and faster acceleration, this will reinforce the instability, possibly triggering "explosive" growth. This can significantly contribute to mechanisms of regulation of P_c discussed earlier in [16].

6. Conclusions and discussion

We have considered a number of possible mechanisms for generation of large scale magnetic field in front of strong astrophysical shocks. All these mechanisms are immediate results of the particle acceleration process. Such generation is necessary to further accelerate particles well beyond the "knee" energy at 10^{15} eV, as is suggested by observations of the CR background spectrum and the wide consensus on their SNR shock origin. The fact that accelerated CRs constitute an ample reservoir for the turbulence which is required to further accelerate them provides a logical basis for our approach. Indeed, the scenario proposed here may be viewed as a self-regulating enhanced acceleration process, which ultimately forces the energetic particle pressure gradient to its marginal point for Drury instability.

A new description of the instability of Alfven wave spectrum to acoustic modulations is given. Along with the Drury instability this is shown to provide an efficient mechanism for transformation of magnetic energy to longer scales. It should be also emphasized that the theoretical analysis of CR acceleration is a challenging problem in plasma wave physics, particle kinetics and shock hydrodynamics. It should, and indeed must, include a self-consistent description of particle transport, wave dynamics, and shock structure.

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