GRAIN ACCELERATION BY MAGNETOHYDRODYNAMIC TURBULENCE: GYRORESONANCE MECHANISM

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ABSTRACT

We discuss a new mechanism of dust acceleration that acts in a turbulent magnetized medium. The magnetohydrodynamic (MHD) turbulence includes both fluid motions and magnetic fluctuations. We show that while the fluid motions bring about grain motions through the drag, the electromagnetic fluctuations can accelerate grains through resonant interactions. In this Letter, we calculate the grain acceleration by the gyroresonance in the cold neutral medium. We consider both incompressible and compressible MHD modes. We show that fast modes dominate the grain acceleration. For the parameters chosen, fast modes render to grains supersonic velocities that may shatter the grains and enable the efficient absorption of heavy elements. Since the grains are preferentially accelerated with large pitch angles, the supersonic grains get aligned with long axes perpendicular to the magnetic field.

Subject headings: acceleration of particles — dust, extinction — ISM: kinematics and dynamics — MHD — turbulence

1. INTRODUCTION

Dust is an important constituent that is essential for the heating and cooling of the interstellar medium (ISM). It interferes with observations in the optical range but provides an insight into star formation activity through far-infrared radiation. It also enables the formation of molecular hydrogen and traces the magnetic field via emission and extinction polarimetry (see Lazarian 2003). The basic properties of dust (optical, alignment, etc.) strongly depend on its size distribution. The latter evolves as the result of grain collisions, whose frequency and consequences (coagulation, cratering, shattering, and vaporization) depend on the relative velocity of grains (see discussions in Draine 1985 and Lazarian & Yan 2002a, 2002b).

All these problems require an understanding of grain motions in the turbulent ISM. Although turbulence has been invoked by a number of authors (see Kusaka, Nakano, & Hayashi 1970, Draine 1985, Ossenkopf 1993, and Weidenschilling & Ruzmaikina 1994) to provide a substantial relative motion of grains, the turbulence they discussed was not magnetized. In a recent paper (Lazarian & Yan 2002a, hereafter LY02), we applied the theory of Alfvénic turbulence (Goldreich & Sridhar 1995, hereafter GS95; see Cho, Lazarian, & Vishniac 2002a for a review) to provide a substantial relative motion of grains, the magnetic field through resonant interactions. In this Letter, we calculate the grain acceleration by the gyroresonance in the cold neutral medium. We consider both incompressible and compressible MHD modes. We show that fast modes dominate the grain acceleration. For the parameters chosen, fast modes render to grains supersonic velocities that may shatter the grains and enable the efficient absorption of heavy elements. Since the grains are preferentially accelerated with large pitch angles, the supersonic grains get aligned with long axes perpendicular to the magnetic field.

Here we account for the acceleration that arises from the resonant interaction of charged grains with MHD turbulence. To describe the turbulence statistics, we use the analytical fits to the statistics of Alfvénic modes obtained in Cho, Lazarian, & Vishniac (2002b, hereafter CLV02) and of compressible modes obtained in Cho & Lazarian (2002, hereafter CL02).

2. ACCELERATION OF GRAINS BY GYRORESONANCE

Turbulent acceleration may be viewed as the acceleration by a spectrum of MHD waves that can be decomposed into incompressible Alfvénic modes and compressible fast and slow modes (see CL02). An important analogy exists between the dynamics of charged grains and the dynamics of comic rays (see Yan & Lazarian 2002, hereafter YL02), and we shall modify the existing machinery used for cosmic rays to describe the charged grain dynamics. The energy exchange involves resonant interactions between the particles and the waves. Specifically, the resonance condition is \( \omega - k_z v_A = n\Omega (n = 0, \pm 1, \pm 2, \ldots) \), where \( \omega \) is the wave frequency, \( k_z \) is the parallel component of wavevector \( k \) along the magnetic field, \( v \) is the particle velocity, \( \mu \) is the cosine of the pitch angle relative to the magnetic field, and \( \Omega = qB/(mc) \) is the Larmor frequency of the particle. The sign of \( n \) denotes the polarization of the wave. The plus sign represents the left-hand polarization, and the minus sign represents the right-hand polarization. Basically, there are two main types of resonant interaction: gyroresonance acceleration and transit acceleration. Transit acceleration \((n = 0)\) requires longitudinal motions and only operates with compressible modes. It happens when \( k_z v_A = \omega \). Since the phase speed is \( V_{\text{fast}} \geq V_\Omega \) for fast waves, where \( V_\Omega \) is the Alfvén speed, it is clear that it can only be applicable to super-Alfvénic grains, which we do not deal with here.

Gyroresonance occurs when the Doppler-shifted frequency of the wave in the grain’s guiding center rest frame \( \omega_{ec} = \omega - k_z v_A \) is a multiple of the particle gyrofrequency and when the rotating direction of the electric wavevector is the same as the direction of the Larmor gyration of the grain. The gyroresonance scatters and accelerates the particles. The efficiency of the two processes for charged grains can be described by the Fokker-Planck coefficients \( D_{\omega c} \) and \( D_{p/p} \), where \( p \) is the particle momentum (see Schlickeiser & Achatz 1993 and YL02). The ratio of the two rates depends on the ratio of the particle velocity, the Alfvén speed, and the pitch angle, \( p^2 D_{\omega c}/D_{p/p} = (\xi \gamma /\gamma_0) + \mu \), where \( \xi = 1 \) for Alfvén waves and \( \xi = k_z /k \) for fast modes. We see that the scattering is less efficient for sub-Alfvénic grains unless most of the particles move parallel to the magnetic field. We shall show later that as the result of acceleration, \( \mu \) will tend to 0. Therefore, in the zeroth-order approximation, we ignore the effect of scattering and assume that the pitch angle cosine \( \mu \) does not change while being accelerated. In this case, the Fokker-Planck equation,
with pitch angle equation (1) we can get the energy gain rate for the grain where is the sound speed. When the grain becomes supersonic (Purcell 1969), the gas drag time by the factor given in Draine & Salpeter (1979). When the grain is larger than the atom-grain cross section. Therefore, in the presence of collisions with ions, the effective drag time decreases by the factor given in Draine & Salpeter (1979). When the grain velocity becomes supersonic (Purcell 1967), the drag time is given by 

\[ t_{\text{drag}} = \frac{t_{\text{drag}}/(0.75 + 0.75e^{-2/5}v^2 - 1/2v^2 + e^{-2/5}v^2)}{c}, \]

where c is the sound speed.

By multiply the Brownian equation above by v and taking the ensemble average, we obtain

\[ m \frac{d\langle v \rangle}{dt} = -\langle \dot{v} \rangle + \langle \dot{v} \rangle, \tag{2} \]

Following an approach similar to that in Melrose (1980), from equation (1) we can get the energy gain rate \( \langle \dot{v} \rangle \) for the grain with pitch angle \( \mu \),

\[ \langle \dot{v} \rangle = \frac{1}{p^2} \frac{\partial}{\partial p} [p^2D_{pp}(\mu)], \tag{3} \]

where the Fokker-Planck coefficient \( D_{pp}(\mu) \) is given below.

Adopting the result from quasi-linear theory (QLT; see Schlickeiser & Achatz 1993), the momentum diffusion coefficient is (see YL02)

\[ D_{pp}(\mu) = \frac{\pi^2(1 - \mu^2)p^3V_i^2}{2v^2} \int \frac{dk}{\tau_{e}^{-1} + (\omega - k_i^\mu - n\Omega)^2} \times \left\{ \begin{array}{l} J_{n+1}(k_{n+1} | \Omega) \left(K_{n+1}(k) + J_{n+1}(k_{n+1} | \Omega) K_{n+1}(k) \right) \\
- J_{n+1}(k_{n+1} | \Omega) \left[e^{i\phi}K_{n+1}(k) + e^{-i\phi}K_{n+1}(k) \right] \end{array} \right\}, \tag{4} \]

where \( \tau_e \) is the nonlinear decorrelation time (and essentially the cascading time of the turbulence), \( k_{n+1} = (k_i^2 + k_{n+1}^2)^{1/2} \) is the perpendicular component of the wavevector, \( v_{n+1} \) is the perpendicular component of the grain velocity, \( \phi = \tan^{-1}(k_i/k_{n+1}) \), \( K_{n+1}(k) \) is the velocity correlation tensor (and will be given in § 3), and \( L, R = (x \pm iy)/\sqrt{2} \) represent left- and right-hand polarization. At the magnetostatic limit \( (\gamma \to \infty) \), the so-called Breit-Wigner-type function transfers into the \( \delta \)-function, i.e., \( \tau = \frac{1}{\gamma} + (\omega - k_i^\mu - n\Omega)^2 \). The magnetostatic limit is correct for fast-moving particles (see YL02), but for sub-Alfvénic grains, we should use equation (4). However, we should not integrate over the whole range of \( k_i \), because the contribution from the large scale is spurious (see YL02). This contribution stems from the fact that in QLT, an unperturbed particle orbit is assumed, which results in the nonconservation of the adiabatic invariant \( \xi = mV_i^2/(2B_0^2) \), where \( B_0 \) is the large-scale magnetic field. Noticing that the adiabatic invariant is conserved only when the electromagnetic field varies on a timescale larger than \( |\Omega|^{-1} \), we truncate our integral range; namely, we integrate from \( k_{\text{res}} \) instead of the injection scale \( L^{-1} \). For Alfvenic turbulence \( \omega = |k_i|V_i \), the resonant scale corresponds to \( |k_i|_{\text{res}} = \Omega/(V_i - \mu R) \). For fast modes in a low-\( \beta \)-medium (where \( \beta = b_{\text{mag}}/B_0^2 = 2/\Omega^2V_i^2 \) is the ratio of gaseous pressure to magnetic pressure), \( \omega = k_i^\mu \), the resonant scale is \( k_{\text{res}} = \Omega/(V_i - \mu R) \). The upper limit of the integral \( k_i \) is set by the dissipation of the MHD turbulence, which varies with the medium.

Integrating from \( k_{\text{res}} \) to \( k_i \), we can obtain from equations (3) and (4) the energy gain rate \( \langle \dot{v} \rangle \) as a function of \( v \) and \( \mu \). Then with \( \langle \dot{v} \rangle \) known, we can estimate the grain acceleration. Solving equation (2) iteratively, we can get the grain velocity as a function of time. We check that the grain velocities converge to a constant value after the drag time. Thus, inserting the acceleration rate by fast modes and Alfven modes into equation (2), we can obtain the final grain velocities as a function of \( \mu \). As \( \langle \dot{v} \rangle \) increases with pitch angle, grains gain the maximum velocities perpendicular to the magnetic field, and therefore the averaged \( \mu \) decreases.

3. MHD TURBULENCE AND ITS TENSOR DESCRIPTION

Unlike hydrodynamic turbulence, Alfvénic turbulence is anisotropic, with eddies elongated along the magnetic field. The Alfvénic turbulence is described by the GS95 model, which postulates that \( k_i V_i \sim k_i V_i \). This may be viewed as the coupling of eddies perpendicular to the magnetic field and wavelike motions parallel to the magnetic field. For the magnetically dominated, so-called low-\( \beta \) plasma, CL02 show that the coupling of Alfvén and compressible modes is weak and that the Alfvén modes and slow modes follow the GS95 spectrum. This is consistent with the analysis of H I velocity statistics (Lazarian & Pogosyan 2000; Stanimirovic & Lazarian 2001). According to CL02, fast modes are isotropic. In what follows, we consider both Alfvén modes and compressible modes in low-\( \beta \) plasma.

Within the random-phase approximation, the velocity correlation tensor in Fourier space is (see Schlickeiser & Achatz 1993)

\[ \langle v_{\alpha}^{(x)}(k, \tau) v_{\beta}^{(x)}(k', \tau' + \tau) \rangle /V_i^2 = \delta(k - k') K_{\alpha\beta}(k) e^{-i\tau}, \]

where \( v_{\alpha\beta} \) is the time-dependent velocity fluctuation in \( k \)-space associated with the turbulence. The velocity correlation tensor for Alfvénic turbulence is (CLV02)

\[ K_{\alpha\beta}(k) = \frac{L^{15/3}}{12\pi} I_2 k_i^{-10/3} \exp\left(-L^{15/3}|k_i V_i|^{-12/3}\right), \]

\[ \tau_i = (L/V_i)(k_i L)^{-12/3} \sim (k_i V_i)^{-1}, \tag{5} \]

where \( I_2 = \delta_{\alpha\beta} - k_i k_j k_i k_j k_j^{-1} \) is a two-dimensional tensor in the \( x\gamma \) plane that is perpendicular to the magnetic field, \( L \) is the injection scale, and \( V \) is the velocity at the injection scale. Velocity fluctuations related to slow modes are subdominant
for magnetically dominated plasmas (CL02), and we do not consider them.

Fast modes are isotropic and have a one-dimensional energy spectrum \( E(k) \propto k^{-3/2} \) (CL02). In the low-\( \beta \) medium, the velocity fluctuations are always perpendicular to \( \mathbf{B} \), for all \( \mathbf{k} \), and the corresponding correlation is (YL02)

\[
K_s(k) = \frac{L^{3/2}}{8\pi} J_y k^{-7/2}, \quad \tau_s = (kL)^{-1/2} V_c^2 / V^2, \tag{6}
\]

where \( J_y = k_y k / k_z^2 \) is also a two-dimensional tensor in the \( x-y \) plane.

4. RESULTS

We consider a typical cold neutral medium (CNM) with \( T = 100 \) K, \( n_a = 30 \) cm\(^{-3} \), and \( B_0 = 6.3 \) \( \mu \)G. Here we only consider large grains (10\(^{-6} \) cm < \( a < 10^{-4} \) cm), which carry the most grains' mass (~80\%) in the ISM. The mean grain charge was obtained from the average electrostatic potential \( \langle U \rangle \) in Weingartner & Draine (2001). MHD turbulence requires that fluid velocities are smaller than the Alfven speed. Therefore, we assume that the injection of energy happens at the scale \( L \), where the equipartition between magnetic and kinetic energies, i.e., \( V = V_c \), is reached. We assume that the velocity dispersion at the scale \( l = 10 \) pc is 5 km s\(^{-1} \) and that the turbulence at large scales proceeds in tenuous warm media with Alfven speeds larger or equal to 5 km s\(^{-1} \). In a partially ionized medium, viscosity caused by neutrals results in decoupling on the characteristic timescale (see LY02),

\[
t_{\text{damp}} \sim v_{\text{th}}^{-1} k^{-2} \sim (l_{p,a})^{-1} k^{-2}, \tag{7}
\]

where \( v_{\text{th}} \) is the kinetic viscosity, \( l_p \) is the neutral mean free path, and \( v_{\text{th}} \) is the thermal velocity of neutrals. Given the parameters above, \( l_p \approx 7 \times 10^{12} \) cm. When its cascading rate \( \tau_{\text{c}} = k / k_{\text{c},T} \approx k^{2/3} L^{-1/3} V_c \) (see eq. [5]) equals the damping rate \( t_{\text{damp}} \), Alfvenic turbulence is assumed to be damped.\(^3\) This defines the cutoff wavenumber of the turbulence \( k_{\text{c},T} = 2.4 \times 10^{-16} \) cm\(^{-1} \) and the timescale \( \tau_c = (kL)^{-1/2} V_c \approx 1.7 \times 10^{10} \) s. Assuming that the grain velocities are smaller than the Alfven speed, we can find that the prerequisite for the gyroresonance \( |k_{\text{c},\nu}| > |k_{\text{c},T}| \) is the same as \( \tau_c > 2 \pi / |\Omega| > \tau_{\text{damp}} \), the condition for effective hydrodynamic drag (see LY02). Thus, we see that Alfven modes cannot accelerate grains (with \( a < 2 \times 10^{-3} \) cm) through gyroresonance unless the velocities of these grains are already super-Alfvenic. The cutoff of fast modes corresponds to the scale for which the cascading timescale \( \tau_c = t_{\text{damp}} \), and this gives the cutoff wavenumber \( k_{\text{c},\nu} = 4.9 \times 10^{-17} \) cm\(^{-1} \). In the present Letter, we consider neutral gas of low ionization, and therefore the damping due to ions, including collisionless damping, is disregarded (cf. YL02). Using the procedure described in § 2, we obtain the grain velocities for different pitch angles. The acceleration is maximal in the direction perpendicular to the magnetic field. Those values are presented in Figure 1. If averaged over \( \mu \), the velocities are smaller by less than 20\%.

In order to compare different processes, we account for the hydrodynamic drag (see LY02 for a discussion of Alfvenic

\(^2\) As pointed out in § 2, the eddy motions happen in the direction perpendicular to the magnetic field, so \( k_z \) is used to calculate \( t_{\text{damp}} \) for Alfven modes.

\(^3\) Thus, we ignore the effect of slowly evolving magnetic structures associated with a recently reported new regime of turbulence below the viscous damping cutoff (Cho, Lazarian, & Vishniac 2002c).

induced drag). The fast modes also cause the relative movement of the grain to the ambient gas by gaseous drag. Unlike gyroresonance, this relative motion arises from the decoupling from the gas. At large scales, grains are coupled with the ambient gas, and the slowing fluctuating gas motions will only cause an overall advection of the grains with the gas (Draine 1985), which we are not interested in. The largest velocity difference occurs on the largest scale where grains are still decoupled. While, in the hydrodynamic case, the decoupling happens on the timescale \( t_{\text{damp}} \) in the MHD case, the grains are constrained by the Larmor gyration unless \( t_{\text{damp}} \). However, the latter condition is true for high-density gas (see Lazarian & Yan 2002b). The velocity fluctuations for fast modes scale as \( v_{\perp} \propto k^{1/4} \propto \omega^{-1/4} \), whereas \( \omega = kV_c \) is the frequency of fast modes (see eq. [6]). In the low-\( \beta \) medium, the velocity fluctuations are perpendicular to the magnetic field. Therefore, grains have velocity dispersions perpendicular to the magnetic field, \( v = V(\tau_c / \tau_{\text{max}})^{3/4} \), where \( \tau_{\text{max}} = L/V \) is the timescale at the injection scale of turbulence. Then we also need to consider the effect of damping. Similar to Alfven modes, the condition for effective hydrodynamic drag \( \tau_i > \tau_{\text{damp}} \), where \( \tau_i = 2 \pi / \omega_i = 5 \times 10^9 \) s is the same as \( k_i > k_{\text{c},\nu} \) for sub-Alfvenic particles, which is the requirement for gyroresonance. Our calculation shows that the corresponding grain size is 4 \( \times 10^{-9} \) cm. For smaller grains, their velocities are reduced, \( v_{\perp} \propto \tau_c / \tau_{\text{damp}} \propto V(\tau_c / \tau_{\text{max}})^{3/4} (\tau_i / \tau_{\text{damp}})^{3/4} \), where \( v_{\perp} \) is the velocity of turbulence at the damping scale (see Lazarian & Yan 2002b).

In Figure 1, we plot the velocity of grains with as a function of grain size since all the mechanisms preferentially accelerate grains in this direction.

How would the results vary as the parameters of the partially ionized medium vary? From equations (2), (3), and (4), we find that the grain velocity is approximately equal to \( v \sim (\bar{\mathbf{v}})_{\text{damp}} / m^{1/2} \sim 2.5 \times 10^{9} L^{-1/4} V_c B^{1/2} (q / a_0)^{0.3} l_{\text{damp}}^{-5/2} \) cm s\(^{-1} \), where \( L_{10} = L / 10 \) pc is the injection scale defined above,
\[ V_s = \frac{V_0}{5} \text{ km s}^{-1}, \quad B_s = B/1 \mu \text{G}, \quad \text{and} \quad q_e = q/1 \text{ electron}. \]

Noticing that the hydrodynamic drag by fast modes decreases with the magnetic field \( (v \propto B^{-1/4} \text{ beyond the damping cutoff, and} \quad v \propto B^{-1} \text{ below the cutoff; see the last paragraph}), \) we see that the relative importance of the gyroresonance and hydrodynamic drag depends on the magnitude of the magnetic field.

It has been shown that the composition of the galactic cosmic ray seems to be better correlated with the volatility of elements (Ellison, Drury, & Meyer 1997). The more refractory elements are systematically overabundant relative to the more volatile ones. This suggests that the material locked in grains must be accelerated more efficiently than gas-phase ions (Epstein 1980; Ellison et al. 1997). The stochastic acceleration of grains, in this case, can act as a preacceleration mechanism.

Grains moving supersonically can also efficiently vacuum-clean heavy elements, as suggested by observations (Wakker & Mathis 2000). Grains can also be aligned if the grains get supersonic (see review by Lazarian 2003). Indeed, the scattering is not efficient for slowly moving grains, so therefore we may ignore the effect of scattering on the angular distribution of the grains. Since the acceleration of grains increases with the pitch angle of the grain (see eqs. [3] and [4]), the supersonic grain motions will result in grain alignment with long axes perpendicular to the magnetic field.

It is believed that silicate grains would not be shattered unless their velocities reach 2.7 \text{ km s}^{-1} (Jones, Tielens, & Hollenbach 1996). Nevertheless, the threshold velocities depend on the velocity of turbulence at the injection scale and on the grain structure, i.e., solid or fluffy. Thus, shattering of the largest grains is possible.

5. SUMMARY

In this Letter, we show that:

1. Fast modes provide the dominant contribution for the acceleration of charged grains. The velocities obtained are sufficiently high to be important for shattering large grains and efficiently absorbing heavy elements from gas. Alfvén modes are not important because of their anisotropy.

2. Depending on the relative importance of the magnetic field, gyroresonance (strong \( B \)) or hydrodynamic drag (weak \( B \)) by fast modes dominates grain acceleration. For small grains, hydrodynamic drag by fast modes is the most important while the gyroresonance is not present as the turbulence at the resonant frequencies gets viscously damped.

3. In the low-\( \beta \) medium, all the mechanisms tend to preferentially accelerate grains in the direction perpendicular to the magnetic field. Among them, gyroresonance with fast modes can render grains with supersonic motions, which can result in grain alignment perpendicular to the magnetic field.

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