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Zonal winds and dipolar vortices in a rotating dusty magnetoplasma

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Abstract. It is shown that the Rossby and dust-Alfvén waves in a rotating dusty magnetoplasma are coupled due to the spatial nonuniformity of the angular rotation velocity of the dust fluid. The nonlinear wave interaction is governed by a pair of equations comprising the evolution of the dust fluid vorticity and the dust-Alfvén wave magnetic field. These nonlinear equations are then used to investigate the generation of zonal winds and the formation of a dipolar vortex. The results are relevant for understanding the origin of the strong turbulence and large scale structures, which are often observed in the planetary magnetospheres.

Contents

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1. Introduction

It is well known that most of the objects in our solar system contain dust and plasmas [1]–[5]. The dust can be charged [6]–[8] due to a variety of processes including the collection of electrons and ions from the background medium, ultra-violet radiation, sputtering, etc. The collective interactions between the charged dust and the plasmas are quite complex, and they lead to many interesting phenomena in the planetary magnetospheres of the Jovian and Saturnian [2, 3] systems as well as in many other astrophysical objects [4], e.g. in photo-evaporated molecular clouds (thought to be 'star nurseries'), galaxies, nebulae, supernova, etc. There are conclusive evidences of various scale size structures (namely braids, kinks, filaments, solitary vortices) in planetary rings [9, 10], astrophysics and cosmology [11]–[13]. However, the origin of those fine and large scale structures is not known, although there have been some attempts [14]–[17] to understand the properties of those fine structures in terms of the coherent vortices that can exist in rotating neutral fluids [14]–[17] and in non-rotating dusty magnetoplasmas [12].

In this paper, we investigate the properties of nonlinearly coupled Rossby–drift-Alfvén (cRdA) waves in a rotating dusty magnetoplasma whose constituents are charged dust, electrons and ions. By employing a multi-fluid description of our dusty plasma system [8], we derive a pair of equations which governs the evolution of the dust fluid vorticity and the dust-Alfvén wave magnetic field, including a nonuniform dust fluid rotation. The equations are then used to investigate the generation of zonal winds and the formation of a dipolar vortex. The latter can be associated with the large scale structures in the Jovian and Saturnian atmospheres, as well as in the atmospheres of Mars and other planets and in some other astrophysical objects (namely giant molecular clouds in the Eagle Nebula, or double galaxies). Zonal winds [18]–[20] and vortices [21] are supposed to play a very important role in regulating turbulent transports in neutral fluids and in astrophysical plasmas.

The present manuscript is organized in the following fashion. In section 2, we derive the governing nonlinear equations for the cRdA waves. The generation of the zonal winds and the formation of a dipolar vortex are considered in section 3. Section 4 contains a summary and possible applications of our work to planetary systems and other astrophysical objects.

2. Derivation of nonlinear equations

We consider a dusty plasma composed of charged dust, electrons, and ions in a uniform magnetic field $\boldsymbol{B}_0 = \hat{z}B_0$, where $\hat{z}$ is the unit vector along the z axis and B_0 is the magnitude of the ambient magnetic field. At equilibrium, we have $Q_d n_d = en_e - Z_i e n_i$, where Q_d is the dust charge ($Q_d = -Z_d e$ for negatively charged dust and $Q_d = Z_d e$ for positively charged dust, Z_d is the number of charges residing on the dust grain surface, and e is the magnitude of the electron charge), n_j is the number density of the particle species j (j equals i for the ions, e for the electrons, and d for the dust), and Z_i is the ion charge number. The dynamics of the cRdA waves in a rotating dusty magnetoplasma is governed by the dust continuity equation

$$\frac{\partial \rho_d}{\partial t} + \nabla \cdot (\rho_d \boldsymbol{v}_d) = 0, \tag{1}$$

the dust momentum equation

$$\rho_d \left(\frac{\partial}{\partial t} + \boldsymbol{v}_d \cdot \nabla \right) \boldsymbol{v}_d = 2\rho_d \boldsymbol{v}_d \times \boldsymbol{\Omega} + n_d Q_d \left(\boldsymbol{E} + \frac{1}{c} \boldsymbol{v}_d \times \boldsymbol{B} \right) - \nabla p_d, \tag{2}$$

the equations of motion for the inertialess electrons and ions, which are, respectively,

$$0 = -n_e e\left(\boldsymbol{E} + \frac{1}{c}\boldsymbol{v}_e \times \boldsymbol{B}\right) - \nabla p_e, \tag{3}$$

and

$$0 = Z_i n_i e\left(\boldsymbol{E} + \frac{1}{c}\boldsymbol{v}_i \times \boldsymbol{B}\right) - \nabla p_i, \tag{4}$$

where $\rho_d = n_d m_d$ is the dust mass density, m_d is the dust mass, $\boldsymbol{v}_j$ is the fluid velocity, p_j is the pressure, $\boldsymbol{\Omega} = \hat{z}\Omega_0(\boldsymbol{r})$, and c is the speed of light in vacuum. The Coriolis force acting on the dust fluid is $2\rho_d \boldsymbol{v}_d \times \boldsymbol{\Omega} \gg 2\rho_i \boldsymbol{v}_i \times \boldsymbol{\Omega}$, where $\rho_i = n_i m_i$ is the ion mass density and m_i is the ion mass.

The electromagnetic fields, $\boldsymbol{E}$ and $\boldsymbol{B}$, are given by the Ampère and Faraday laws, which are, respectively,

$$\nabla \times \boldsymbol{B} = \frac{4\pi}{c}\boldsymbol{j}, \tag{5}$$

and

$$\frac{\partial \boldsymbol{B}}{\partial t} = -c\nabla \times \boldsymbol{E}, \tag{6}$$

where $\boldsymbol{j} = Z_i e n_i \boldsymbol{v}_i - e n_e \boldsymbol{v}_e + Q_d n_d \boldsymbol{v}_d$ is the plasma current density. We have neglected the displacement current in equation (5) since the phase speed of the cRdA waves is much smaller than c.

Combining equations (3), (4), and (5), we obtain

$$\boldsymbol{E} = -\frac{\boldsymbol{v}_d \times \boldsymbol{B}}{c} - \frac{\nabla\left(P_{ei} + B^2/8\pi\right)}{Q_d n_d} + \frac{(\boldsymbol{B}\cdot\nabla)\boldsymbol{B}}{4\pi Q_d n_d}, \tag{7}$$

where $P_{ei} = p_e + p_i$. Accordingly, equations (2) and (6) can be written as

$$\rho_d\left(\frac{\partial}{\partial t} + \boldsymbol{v}_d \cdot \nabla\right)\boldsymbol{v}_d = 2\rho_d \boldsymbol{v}_d \times \boldsymbol{\Omega} - \nabla\left(P + \frac{B^2}{8\pi}\right) + \frac{(\boldsymbol{B}\cdot\nabla)\boldsymbol{B}}{4\pi}, \tag{8}$$

and

$$\frac{\partial \boldsymbol{B}}{\partial t} = \nabla \times (\boldsymbol{v}_d \times \boldsymbol{B}) - \frac{c}{4\pi}\nabla \times \left[\frac{(\boldsymbol{B}\cdot\nabla)\boldsymbol{B}}{Q_d n_d}\right], \tag{9}$$

where $P = P_{ei} + p_d$. Equations (8) and (9) are the magnetohydrodynamic equations for a rotating dusty magnetoplasma provided that the equation of state for the total plasma pressure is known. Its specific form is, however, not needed below as the curl of the gradient of the pressure is zero.

We now consider the two-dimensional dust fluid motions in a magnetoplasma with $2\boldsymbol{\Omega} = 2\Omega_0\hat{z}\sin\lambda \equiv \hat{z}f$, where λ is the latitude of the site counted from the equator on the planet, the x, y, and z axes are directed eastwards, northwards, and upwards, respectively, f is the Coriolis frequency, $f = f_0 + \beta y$, where $f_0 = 2\Omega_0 \sin\lambda_0$ and $\beta = (2\Omega_0/r_0)\cos\lambda_0 > 0$. Here, within the β-plane approximation, λ_0 is the latitude of the site, r_0 is the distance from the center of the planet, and $y = r_0(\lambda - \lambda_0)$. For two-dimensional dust motions, we can set $v_{dz} = 0$ and $B_z = 0$ and take $\nabla \cdot \boldsymbol{v}_{d\perp} \approx 0$ and $\nabla \cdot \boldsymbol{B}_\perp = -\partial B_z/\partial z \approx 0$, where v_{dz} and B_z are the components of the dust fluid velocity and the cRdA wave magnetic field along the z axis, and

$\boldsymbol{v}_{d\perp}(x,y,t)$ and $\boldsymbol{B}_\perp(x,y,t) = \boldsymbol{B}(x,y,t) - \hat{z}B_0$ are the transverse (to $\hat{z}$) components of the dust fluid velocity and the cRdA wave magnetic field, respectively. We then write

$$\boldsymbol{v}_{d\perp} = \hat{z}\times\nabla_\perp\phi \quad \text{and} \quad \boldsymbol{B}_\perp = -\hat{z}\times\nabla_\perp\psi, \tag{10}$$

where ϕ and ψ are stream functions and $\nabla_\perp = \hat{x}(\partial/\partial x) + \hat{y}\partial/\partial y$. Such choices for the dust fluid velocity and magnetic field perturbations correspond to a dust fluid vorticity $\boldsymbol{\Omega}_d = \hat{z}\nabla_\perp^2\phi$ and a parallel current density $\propto \hat{z}\nabla_\perp^2\psi$.

Taking the curl of equation (8) we have for constant ρ_d and Q_d the dust vorticity equation

$$\left(\frac{\partial}{\partial t} + \hat{z}\times\nabla_\perp\phi\cdot\nabla\right)\nabla_\perp^2\phi + \beta\frac{\partial\phi}{\partial x} = -\frac{v_{\mathrm{A}d}}{\sqrt{4\pi\rho_d}}\left(\frac{\partial}{\partial z} - \frac{1}{B_0}\hat{z}\times\nabla_\perp\psi\cdot\nabla\right)\nabla_\perp^2\psi, \tag{11}$$

where $v_{\mathrm{A}d} = B_0/\sqrt{4\pi\rho_d}$ is the dust-Alfvén speed and $\nabla_\perp^2 = \partial^2/\partial x^2 + \partial^2/\partial y^2$ is the transverse Laplacian. On the other hand, equation (9) gives the evolution equation for the parallel current density

$$\left(\frac{\partial}{\partial t} + \hat{z}\times\nabla_\perp\phi\cdot\nabla\right)\psi = -B_0\frac{\partial\phi}{\partial z}. \tag{12}$$

The coupled equations (11) and (12) are the desired set for the cRdA waves in an astrophysical dusty fluid that is rotating with a nonuniform speed. They conserve the energy integral

$$W = \int \mathrm{d}x\mathrm{d}y\left[\frac{1}{2}(\nabla_\perp\phi)^2 + \frac{1}{8\pi\rho_d}(\nabla_\perp\psi)^2\right]. \tag{13}$$

In the linear approximation, we can neglect the nonlinear Poisson bracket terms in equations (11) and (12) and Fourier transform the resultant equations by considering the stream functions to be proportional to $\exp(-\mathrm{i}\omega t + \mathrm{i}\boldsymbol{k}\cdot\boldsymbol{r})$, where ω and $\boldsymbol{k}$ are the frequency and the wavevector, respectively. Combining the Fourier transformed equations, we then obtain the linear dispersion relation

$$\omega^2 + \omega\omega_{\mathrm{R}} - \omega_{\mathrm{A}d}^2 = 0, \tag{14}$$

which exhibits a linear coupling between the Rossby and dust-Alfvén waves. The Rossby and dust-Alfvén wave frequencies are denoted by, respectively,

$$\omega_{\mathrm{R}} = \frac{2\Omega_0\cos\lambda_0}{r_0}\frac{k_x}{k_x^2 + k_y^2}, \tag{15}$$

and

$$\omega_{\mathrm{A}d} = k_z v_{\mathrm{A}d}. \tag{16}$$

For $k_z = 0$, the Rossby waves ($\omega = -\omega_{\mathrm{R}}$) are degenerated.

3. Zonal winds and vortices

In this section, we present analytical studies of the zonal wind excitation by the cRdA waves, as well as of the formation of a dipolar vortex in a rotating dusty magnetoplasma. To study the zonal wind excitation, we write $\phi = \Phi + \varphi$ and $\psi = \Psi$ in equations (11) and (12), where Φ and Ψ are the stream functions associated with the cRdA, and φ is the stream function of the zonal winds which have a variation only along the y axis. We then find for the cRdA waves the equations

$$\frac{\partial}{\partial t}\nabla_\perp^2\Phi + \beta\frac{\partial\Phi}{\partial x} + \left(\hat{z}\times\nabla_\perp\varphi\cdot\nabla\right)\nabla_\perp^2\Phi + \left(\hat{z}\times\nabla_\perp\Phi\cdot\nabla\right)\nabla_\perp^2\varphi + \frac{v_{\mathrm{Ad}}}{\sqrt{4\pi\rho_d}}\frac{\partial}{\partial z}\nabla_\perp^2\Psi = 0, \tag{17}$$

and

$$\frac{\partial\Psi}{\partial t} + B_0\frac{\partial\Phi}{\partial z} + \hat{z}\times\nabla_\perp\varphi\cdot\nabla\Psi + \hat{z}\times\nabla_\perp\Psi\cdot\nabla\varphi = 0. \tag{18}$$

In a similar manner, the equation for the zonal winds is found to be

$$\frac{\partial}{\partial t}\nabla_\perp^2\varphi + \langle\left(\hat{z}\times\nabla_\perp\Phi\cdot\nabla\right)\nabla_\perp^2\Phi - \frac{v_{\mathrm{Ad}}}{B_0\sqrt{4\pi\rho_d}}(\hat{z}\times\nabla_\perp\Psi\cdot\nabla)\nabla_\perp^2\Psi\rangle = 0, \tag{19}$$

where the second and third terms in the left-hand side of equation (19) represent the stresses associated with the cRdA waves. The angular bracket in (19) denotes averaging over the cRdA wave period. In equation (19), Ψ is determined from $\partial\Psi/\partial t + B_0\partial\Phi/\partial z = 0$. Equations (18) and (19) are a relevant set for studying the excitation of zonal winds in a rotating dusty magnetoplasma.

We now consider the parametric excitation of the zonal winds by cRdA waves, given by equation (14). The nonlinear interaction between a cRdA pump (ω_0, $\boldsymbol{k}_0$) and the zonal winds (ω, $\boldsymbol{k}$) excite upper and lower sidebands ($\omega_\pm$, $\boldsymbol{k}_\pm$). Thus, we decompose the stream functions as

$$\Phi = \Phi_{0+}\exp(-\mathrm{i}\omega_0 t + \mathrm{i}\boldsymbol{k}_0\cdot\boldsymbol{r}) + \Phi_{0-}\exp(\mathrm{i}\omega_0 t - \mathrm{i}\boldsymbol{k}_0\cdot\boldsymbol{r}) + \sum_{+,-}\Phi_\pm\exp(-\mathrm{i}\omega_\pm t + \mathrm{i}\boldsymbol{k}_\pm\cdot\boldsymbol{r}), \tag{20}$$

and

$$\Psi = \Psi_{0+}\exp(-\mathrm{i}\omega_0 t + \mathrm{i}\boldsymbol{k}_0\cdot\boldsymbol{r}) + \Psi_{0-}\exp(\mathrm{i}\omega_0 t - \mathrm{i}\boldsymbol{k}_0\cdot\boldsymbol{r}) + \sum_{+,-}\Psi_\pm\exp(-\mathrm{i}\omega_\pm t + \mathrm{i}\boldsymbol{k}_\pm\cdot\boldsymbol{r}), \tag{21}$$

and

$$\varphi = \tilde{\varphi}\exp(-\mathrm{i}\omega t + \mathrm{i}\boldsymbol{k}\cdot\boldsymbol{r}), \tag{22}$$

where $\omega_\pm = \omega \pm \omega_0$ and $\boldsymbol{k}_\pm = \boldsymbol{k} \pm \boldsymbol{k}_0$ are the frequencies and wave vectors of the sidebands, and the superscript 0 ($\pm$) stands for the pump (sidebands). Such a Fourier decomposition of modes, satisfying the energy and momentum conservation relations, is a unique characteristic of the mode coupling processes within the framework of weak turbulence theory in astrophysics [4].

Inserting (20)–(22) into (17)–(19) and Fourier analysing, we obtain

$$\Omega_\pm^2\Phi_\pm = \pm\mathrm{i}\omega_\pm\frac{\hat{z}\times\boldsymbol{k}_0\cdot\boldsymbol{k}}{k_{\perp\pm}^2}(k_\perp^2 - k_{\perp0}^2)\Phi_{0\pm}\tilde{\varphi}. \tag{23}$$

and

$$\omega\tilde{\varphi} = \mathrm{i}\left(1 - \frac{k_{z0}^2 v_{\mathrm{Ad}}^2}{\omega_0^2}\right)\frac{\hat{z}\times\boldsymbol{k}_0\cdot\boldsymbol{k}}{k_\perp^2}\left(K_-^2\Phi_{0+}\Phi_- - K_+^2\Phi_{0-}\Phi_+\right), \tag{24}$$

where $\Omega_\pm^2 = \omega_\pm^2 + \omega_\pm\omega_{\mathrm{R}\pm} - k_{z\pm}^2 v_{\mathrm{Ad}}^2$, $\omega_{\mathrm{R}\pm} = \beta k_{x\pm}/k_{\perp\pm}^2$, $K_\pm^2 = k_{\perp\pm}^2 - k_{\perp0}^2$, $\omega_0^2 + \omega_0\omega_{R0} - k_{z0}^2 v_{Ad}^2 = 0$, and $\omega_{\mathrm{R}0} = \beta k_{x0}/k_{\perp0}^2$.

Eliminating $\Phi_\pm$ from equation (23) by using equation (24), we have the nonlinear dispersion relation

$$\omega \approx -\frac{\omega_{R0}}{\omega_0}\frac{|\hat{z}\times \boldsymbol{k}_0\cdot \boldsymbol{k}|^2}{k_\perp^2}(k_\perp^2 - k_{\perp 0}^2)\sum_{+,-}\frac{K_\pm^2\omega_\pm|\Phi_0|^2}{k_{\perp\pm}^2\Omega_\pm^2}, \tag{25}$$

where $|\Phi_0|^2 = \Phi_{0+}\Phi_{0-}$. Equation (25) admits instability for $k_{\perp 0} \gg k_\perp$ and $\boldsymbol{k}_0\cdot\boldsymbol{k}_\perp > 0$. The growth rate is

$$\gamma \approx \left(\frac{\omega_{R0}}{\omega_0}\right)^{1/2}\left|\frac{\hat{z}\times\boldsymbol{k}_0\cdot\boldsymbol{k}}{k_\perp}\right| |\boldsymbol{k}_0\cdot\boldsymbol{k}_\perp|^{1/2}\,|\Phi_0|. \tag{26}$$

Next, we present the stationary solutions of equations (11) and (12) in a new frame $\xi = x+\alpha z-ut$, where α and u are constants. Supposing that ϕ and ψ are functions of ξ and y only, we obtain then

$$\mathcal{L}\left(\nabla_\perp^2\phi - \frac{\beta}{u}\phi\right) - \frac{\alpha v_{Ad}}{u\sqrt{4\pi\rho_d}}\left(\frac{\partial}{\partial\xi} - \frac{1}{\alpha B_0}\hat{z}\times\nabla_\perp\psi\cdot\nabla\right)\nabla_\perp^2\psi = 0, \tag{27}$$

and

$$\mathcal{L}\left(\psi - \frac{\alpha B_0}{u}\phi\right) = 0, \tag{28}$$

where $\mathcal{L} = (\partial/\partial\xi) - u^{-1}\hat{z}\times\nabla_\perp\phi\cdot\nabla$. A possible solution of equation (28) is

$$\psi = \frac{\alpha B_0}{u}\phi. \tag{29}$$

Hence, equation (27) can be expressed as

$$\mathcal{L}\left[\left(1 - \frac{\alpha^2 v_{Ad}^2}{u^2}\right)\nabla_\perp^2\phi - \frac{\beta}{u}\phi\right] = 0. \tag{30}$$

Equation (30) is satisfied by the ansatz

$$\nabla_\perp^2\phi = C_1\phi - C_2 y, \tag{31}$$

where the constants C_1 and C_2 are related by

$$\left(1 - \frac{\alpha^2 v_{Ad}^2}{u^2}\right)\left(C_1 + \frac{C_2}{u}\right) - \frac{\beta}{u} = 0. \tag{32}$$

The localized dipolar vortex solution of equation (31) is well known [21, 22]. It is given by

$$\phi^{o}(r,\theta) = \phi_o K_1(pr)\sin\theta, \tag{33}$$

in the outer region, $r > a$, and

$$\phi^{i}(r,\theta) = \left[\phi_i J_1(qr) - \frac{p^2+q^2}{q^2}ur\right]\sin\theta, \tag{34}$$

in the inner region, $r < a$, where $r = (\xi^2+y^2)^{1/2}$, $\sin\theta = y/r$, a is the vortex radius, ϕ_o and ϕ_i are arbitrary constants, K_1 is the McDonald function and J_1 is the Bessel function of the first kind, and $p^2 = \beta u/(u^2 - \alpha^2 v_{Ad}^2) > 0$. Through the boundary conditions (continuity of ϕ and $\partial\phi/\partial r$) at the vortex interface we obtain the constants $\phi_o = p^2 ua/q^2 J_1(qa)$ and $\phi_i = ua/K_1(pa)$. For a given value of p, we determine q from the transcendental equation

$$\frac{J_2(qa)}{qaJ_1(qa)} = -\frac{K_2(pa)}{paK_1(pa)}. \tag{35}$$

Thus, a dipolar vortex involving cRdA waves is formed when $u > \alpha v_{\mathrm{A}d}$. For $\alpha = 0$ the translational speed of the dipolar vortex is arbitrary. In fact, the dust fluid vorticity is decoupled from the dust-Alfvén wave magnetic field if $\partial/\partial z = 0$. Here, the dynamics of the flute-like Rossby waves is governed by

$$\left(\frac{\partial}{\partial t} + \hat{z} \times \nabla_\perp \phi \cdot \nabla\right) \nabla_\perp^2 \phi + \beta \frac{\partial \phi}{\partial x} = 0, \tag{36}$$

which is the modified Navier–Stokes equation, admitting dual cascade. The latter ensures the formation of a vortex.

4. Summary

In summary, we have presented an investigation of the linear and nonlinear cRdA waves in a rotating dusty magnetoplasma. By employing a multi-fluid approach and Faraday's and Ampère's laws, we have derived a pair of equations which shows a coupling between the dust fluid vorticity and the magnetic field of the dust-Alfvén waves. The nonlinearly coupled Rossby–dust-Alfvén waves exchange momentum and energy among themselves and can generate zonal winds in a rotating magnetized dust fluid. Thus, we have found a novel nonlinear mechanism via which the zonal winds are excited at the expense of the free energy of finite amplitude short scale cRdA waves. Furthermore, the nonlinear self-interaction between the cRdA waves can produce a coherent dipolar vortex of an arbitrary size, whose speed is larger than $\alpha v_{\mathrm{A}d}$. The results of the present investigation thus offer efficient mechanisms for the excitation of zonal winds and for the formation of vortical structures, which can control the transport processes in planetary systems (namely the atmospheres of Jupiter, Saturn, and Mars) and in other contexts [11], [23]–[25] (namely in giant molecular clouds in our galaxy and in rotating interstellar clouds). Finally, we stress that our results may also be useful in understanding the jet streams, the zonal winds and the vortex motions in the Earth's mesosphere where charged dust particles are inherently present and the dust fluid is rotating. It is hoped that the forthcoming data from the Cassini mission and rocket campaigns will offer more information regarding the zonal winds and vortex motions in the atmosphere of Saturn and in the polar mesosphere of the Earth.

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