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Magnetisation of partially ionised dusty disks

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Abstract. The self-magnetisation of circumstellar disks is considered within an appropriate multifluid description. These disks are composed of ionised and neutral gas as well as of a charged dust component. The most important equation in this context is the general Ohm's law that includes a magnetic field generation term due to relative dust–neutral fluid velocities. We show that circumstellar disks can carry their own significant magnetic fields. As long as the stellar gravitation sustains the accretion flow, the self-magnetisation of the disk does not saturate until the field strength reaches its local equipartition value. The magnetic field generation process is illustrated by idealised multifluid simulations that are not restricted to a kinematic description, but model the process in a self-consistent way.

1. Introduction

A characteristic feature of partially ionised plasmas is the possibility that magnetic fields can be generated without any seed fields. As shown by Huba and Feder (1993) in a fundamental paper, a relative plasma–neutral gas flow with non-vanishing vorticity results inevitably in the generation of magnetic fields. This effect is described by the generalised Ohm's law of a multifluid description. It is to be expected that this magnetisation mechanism works even more efficiently in a dusty plasma (Birk et al. 1996; Shukla et al. 1998), where the inertia effects and collisional momentum transfers are dominated by the heavy dust grains. Investigations of this kind of self-magnetisation of partially ionised dusty plasmas have to be based on the detailed self-consistent balance equations of the different species (Ciolek and Mouschovias 1993; Birk et al. 1996; Kopp et al. 1997; Schröer and Kopp 2000). In this contribution, we consider as an important application the magnetisation of circumstellar accretion disks. The kind of interaction of stellar objects with their associated disks is crucial for dynamic processes such as winds and bipolar flows (Königl 1994; Paatz and Camenzind 1996) and, in particular, magnetic flares (André 1996; Grosso et al. 1997; Stelzer and Neuhäuser 2000). The interaction with a disk that does not support its own magnetic field is different from the interaction with a magnetised one. The question whether or not accretion disks around stellar objects are considered to be magnetised is still not answered conclusively. In the particular case of T Tauri stars, we come to the conclusion that the circumstellar disks should be fully magnetised. After the stellar dynamo has produced a significant stellar bipolar magnetic field, the Keplerian motion is faster than corotation (e.g. Li 1996). Beyond the corotation radius, the disk particles move on Keplerian orbits. Perturbation of these orbits can lead to a starward flow of the particles, if the centrifugal forces no longer balances the gravitational force exactly. The charged particles continue to move inwards until the stellar magnetic field is strong enough to force them to corotate. On the other hand, the neutral particles are not forced to corotate, since they are not coupled to magnetic field lines directly but rather via dust–neutral and ion–neutral momentum transfer and continue to move starwards. The resulting macroscopic sheared relative dust–neutral drift motions are responsible for the generation of magnetic fields. The ultimate cause of the inward motion is the stellar gravitational potential.

2. Basic equations

The macroscopic dynamics of partially ionised dusty plasmas are governed by multifluid balance equations for the different species. We are dealing with isothermal low-frequency (with respect to the dust gyrofrequency) dynamics in partially ionised dusty plasmas that are in ionisation equilibrium. Thus we have to take into account that dust dynamics plays an important role (Pilipp et al. 1987; Ciolek and Mouschovias 1993; Birk et al. 1996; Shukla et al. 1998).

We assume quasineutrality and neglect collisions of electrons with ions and neutrals as compared with other collisional processes. For the dust grains, we assume a homogeneous and constant negative electrical charge, which can be justified in a simple scenario (Bliokh et al. 1995) where the dust is charged due to electron/ion currents that result in a stationary dust charge after $t_{\text{charge}} \approx 4\pi T a e^2 n_e v_e$ (T, a, n, and v_e denote the temperature, the dust grain radius, the electron density, and the electron thermal velocity, respectively). The entire physics of dust charging is very involved (de Angelis 1992; Lafon 1996), but for our purposes it is not necessary to overly dwell on the microphysics, since the charging is much faster than the shortest time scales of the fluid dynamics given by the inverse dust gyrofrequency.

The continuity equations for the three dynamically relevant fluids (ions, dust, and neutrals) are

$$\frac{\partial \rho_\alpha}{\partial t} + \boldsymbol{\nabla} \cdot (\rho_\alpha \mathbf{v}_\alpha) = 0, \tag{2.1}$$

where ρ and $\mathbf{v}$ denote the mass densities and the bulk velocities, and the index α represents the ion (i), dust (d), and neutral (n) species. It should be noted that, in principle, one can compute the electron density by means of the quasineutrality condition. A more detailed multifluid description is given by Kopp et al. (1997). The inertia forces of the ions are assumed to be small in comparison with that of the dust grains and neutrals. The inertialess ion momentum transfer equation reads

$$0 = n_i e \left(\mathbf{E} + \frac{1}{c} \mathbf{v}_i \times \mathbf{B} \right) - \boldsymbol{\nabla} p_i - \rho_i \nu_{id} (\mathbf{v}_i - \mathbf{v}_d) - \rho_i \nu_{in} (\mathbf{v}_i - \mathbf{v}_n) \tag{2.2}$$

and the dust momentum transfer equation is

$$\begin{aligned} \frac{\partial (\rho_d \mathbf{v}_d)}{\partial t} = & -\boldsymbol{\nabla} \cdot (\rho_d \mathbf{v}_d \mathbf{v}_d) - n_d z_d e \left(\mathbf{E} + \frac{1}{c} \mathbf{v}_d \times \mathbf{B} \right) - \boldsymbol{\nabla} p_d + \rho_d \mathbf{g} \\ & -\rho_d \nu_{di} (\mathbf{v}_d - \mathbf{v}_i) - \rho_d \nu_{dn} (\mathbf{v}_d - \mathbf{v}_n), \end{aligned} \tag{2.3}$$

where p, n, e, z_d, c, $\mathbf{E}$, $\mathbf{B}$, and $\mathbf{g}$ are, respectively, the thermal pressures and particle densities of the components, the elementary charge, the dust charge number, the velocity of light, the electric field, the magnetic field, and the gravitational acceleration. The elastic collision frequencies $\nu_{\alpha\beta}$ satisfy the relation $n_\alpha \nu_{\alpha\beta} = n_\beta \nu_{\beta\alpha}$. The set of momentum equations is completed by the momentum transfer equation for the neutrals:

$$\frac{\partial(\rho_n \mathbf{v}_n)}{\partial t} = -\boldsymbol{\nabla}\cdot(\rho_n \mathbf{v}_n \mathbf{v}_n) - \boldsymbol{\nabla} p_n + \rho_n \mathbf{g} - \rho_n \nu_{ni}(\mathbf{v}_n - \mathbf{v}_i) - \rho_n \nu_{nd}(\mathbf{v}_n - \mathbf{v}_d). \quad (2.4)$$

Substituting for $\mathbf{E}$ from Faraday's law

$$\frac{\partial \mathbf{B}}{\partial t} = -c\boldsymbol{\nabla}\times\mathbf{E}, \quad (2.5)$$

and making use of Ampère's law

$$\mathbf{j} = \frac{c}{4\pi}\boldsymbol{\nabla}\times\mathbf{B} = e(n_i \mathbf{v}_i - n_e \mathbf{v}_e - n_d z_d \mathbf{v}_d), \quad (2.6)$$

we obtain from (2.2) the induction equation that governs the dynamical evolution of the magnetic field:

$$\begin{aligned}\frac{\partial \mathbf{B}}{\partial t} = &-\frac{m_i c}{e}\boldsymbol{\nabla}\times\left(\frac{\boldsymbol{\nabla} p_i}{\rho_i}\right) + \frac{m_i}{m_d} z_d \boldsymbol{\nabla}\times\left(\frac{\rho_d}{\rho_i}\mathbf{v}_d\times\mathbf{B}\right)\\ &+\frac{m_i c}{4\pi e}\boldsymbol{\nabla}\times\left(\frac{\boldsymbol{\nabla}\times\mathbf{B}}{\rho_i}\times\mathbf{B}\right) - \eta\nabla^2\mathbf{B} + \mathbf{R}\\ &-\frac{m_i c}{e}\boldsymbol{\nabla}\times\left\{\frac{n_d}{n_i}\left[\left(z_d - \frac{n_i}{n_d}\right)\nu_{id} + z_d\nu_{in}\right]\mathbf{v}_d - \nu_{in}\mathbf{v}_n\right\}. \end{aligned} \quad (2.7)$$

In (2.7), the first four terms on the right-hand side are well known from pure magnetohydrodynamics. They are the pressure term, the convective term, the Hall term, and the magnetic field diffusion term with the collisional (constant) electrical diffusivity $\eta = m_i c^2(\nu_{id} + \nu_{in})/4\pi n_i e^2$. The following term $\mathbf{R}$ (for details, see Kopp et al. 1997) represents the electron contribution to the electric current. This is negligible if the electrons are almost completely attached to the dust grains.

The final term represents the magnetic field self-generation. It is this term that we shall focus on. In an initially magnetic field-free plasma, this term can give rise to the formation of magnetic fields, similar to a battery due to a thermal pressure with non-parallel isochores and isothermals (cf. the first term on the right-hand side of (2.7)) known as Biermann's battery mechanism (Biermann 1950). Huba and Fedder (1993) discussed at length (with applications to planetary ionospheres and cometary boundary layers) how such a magnetic field self-generation caused by relative flows with non-vanishing vorticities can operate in partially ionised dust-free plasmas. The extension to dusty plasmas is straightforward. The key condition for the considered mechanism to work is a macroscopic sheared relative drift between the dust and the neutral component. Such a macroscopic drift seems to be unavoidable in the context of circumstellar disks (see the next section).

It is worth mentioning that such a magnetic field generation term in a non-static dusty plasma would also be present in the case of a positively (which would be dominant, if photoionisation was the main charging process) or mixed charged dust component. In particular, the contribution $\sim \nu_{in}$ would be exactly the same for positively charged dust grains.

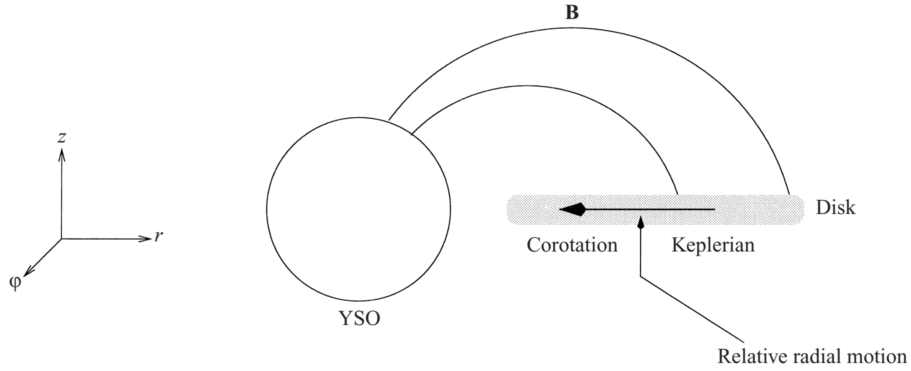

Figure 1. Accretion disk around a young stellar object (YSO). Within the disk, a relative motion of the charged dust and the neutral fluid components due to different coupling properties to the stellar magnetic field can occur. For the coordinates chosen, the flow is in the r direction, the z direction is perpendicular to the disk, and the φ axis points in the azimuthal direction.

3. Application: magnetisation of circumstellar disks

An important astrophysical application of the discussed scenario is to the disks around young stellar objects (YSO). These are believed to be the result of gravitational instabilities of rotating molecular clouds. The outer parts of these clouds should form accretion disks due to an excess of angular momentum (Bodenheimer 1993). Perturbations of the Keplerian orbits of the disk particles should result in a separation of the charged and neutral fluid particles: whereas the charged particles, in particular the massive dynamical dust particles, are forced to corotate due to their coupling to the stellar magnetic field lines, an efficient mass transport of the neutral fluid perpendicular to the stellar magnetic dipole field will occur, since the neutrals are braked by the slower corotating charged particles and thus become too slow for stable Keplerian motion. Consequently, a radial, non-saturating relative macroscopic drift motion of the neutral and the dust fluid components will set in (Fig. 1). The effect of this relative motion, which is ultimately driven by the external stellar gravitational field, on the temporal evolution of magnetic fields is governed by the last term of (2.7):

$$\frac{\partial \mathbf{B}_{\text{gen}}}{\partial t} \sim -\frac{m_i c}{e} \boldsymbol{\nabla} \times \left\{ \frac{n_d}{n_i} \left[\left(z_d - \frac{n_i}{n_d} \right) \nu_{id} + z_d \nu_{in} \right] \mathbf{v}_d - \nu_{in} \mathbf{v}_n \right\}. \tag{3.1}$$

Note that the amplitude of the generated magnetic field does not depend on the actual values of the particle densities of the ions and the dust grains, but only on the ratio n_d/n_i. Obviously, if the quasineutrality condition reads $n_i \approx z_d n_d$ (which means that the electron density is effectively depleted), the generation term can simply be estimated as a function of the radial dust–neutral shearing velocity, with an amplitude measured by the ion–neutral collision frequency.

In a cylindrical geometry (Fig. 1), any inhomogeneity in the z direction results in a relative sheared flow, which gives rise to the generation of a poloidal disk magnetic field (cf. (3.1)):

$$\frac{\partial \mathbf{B}_{\text{gen}}}{\partial t} = \dot{B}_\varphi \mathbf{e}_\varphi \sim \frac{\partial}{\partial z} (\mathbf{v}_d - \mathbf{v}_n)_r \mathbf{e}_\varphi. \tag{3.2}$$

The ion–dust and ion–neutral elastic collision frequencies can be calculated from the appropriate Landau collision integrals (Benkadda et al. 1996; Huba 1998). The ion–neutral collision frequency can be estimated as $\nu_{in} \approx 5 \times 10^{-15} n_n \sqrt{kT_i/m_i}$, where k is the Boltzmann constant and T_i is the ion temperature, and the ion–dust collision frequency is

$$\nu_{id} \approx \frac{4\sqrt{2\pi} n_d z_d^2 e^4 \ln(\lambda_D/a)}{3\sqrt{m_i (kT)^3}},$$

where $\lambda_D = \sqrt{kT_d/4\pi n_d z_d^2 e^2}$ is the dust Debye screening length and T_d is the dust temperature.

This leads to the following expression for the generation of the poloidal magnetic field in the accretion disk:

$$\dot{B}_\varphi \approx \frac{1}{L_z}\left\{ \left[10^{-5}\left(\frac{n_d z_d}{n_i} - 1\right)\frac{n_d z_d^2}{T_i^{3/2}} \ln\left(\frac{\lambda_D}{a}\right) + 5 \times 10^{-15} n_n \sqrt{T_i}\frac{n_d z_d}{n_i}\right] (\mathbf{v}_d)_r \right.$$
$$\left. -5 \times 10^{-15}\, n_n \sqrt{T_i} (\mathbf{v}_n)_r \right\}, \tag{3.3}$$

where L_z denotes the shear length, i.e. the length scale of the inhomogeneity of the relative dust-neutral flow.

For quantitative estimates, we need to specify a parameter set. As an example we choose physical quantities typical for circumstellar disks around T Tauri stars (Montmerle 1991; Königl 1994). The neutral-gas temperature and the mass of the dust grains have no influence on our results. Moreover, the results do not depend sensitively on the grain size, chosen as $a = 10^{-8}$ cm, since it only enters (as does the dust temperature) in the logarithmic term for the ion–dust collision frequency. We obtain a lower limit for the poloidal magnetic field in the accretion disk if we assume $n_i \approx n_d z_d$ (note that in this case of almost total electron depletion, the term proportional to ν_{in} is the only finite generation term in (2.7)):

$$\dot{B}_\varphi \approx 10^{-8}\,\mathrm{G\,s^{-1}} \left[\frac{n_n}{10^{12}\,\mathrm{cm^{-3}}}\right] \sqrt{\left[\frac{T_i = T_d}{5 \times 10^2\,\mathrm{K}}\right]} \left[\frac{L_z}{R_{\mathrm{Sun}}}\right]^{-1} \left[\frac{v_d - v_n}{v_{ff}}\right] \left[\frac{z_d}{1}\right]. \tag{3.4}$$

The choice of the shear length as one solar radius ($L_z = R_{\mathrm{Sun}}$) should mean again that we underestimate B_{gen}. The charge number was chosen as $z_d = 1$ (implying $n_d/n_i = 1$), which also is a lower limit, and the dust mass as $m_d = 1000 m_i$. The relative dust–neutral motion is measured in units of the free-fall velocity $v_{ff} = \sqrt{2G\,M_{\mathrm{Sun}}/D}$ (where G is the gravitational constant and M_{Sun} is the solar mass) given that the free-fall time is the shortest characteristic time scale of the system. We choose $D = 10$ AU for the distance from the central object. It should be noted that, for a slight deviation from $n_i = n_d z_d$, ion–dust collisions also contribute to the generation of magnetic fields, with an even higher efficiency.

In the case considered, the magnetic field self-generation is limited by the free-fall time $t_{ff} = 2\pi D/v_{ff} = 7 \times 10^8$ s. This is the shortest time scale involved, and results in a disk magnetic field of the order of $B_\varphi \approx 10$ G. The other extreme estimate considers the accretion time. With the α-prescription (Shakura and Sunyaev 1973; von Rekowski and Fröhlich 1997) for the turbulent viscosity $\zeta_{\mathrm{turb}} = \alpha c_s L_z$ (where $\alpha = 1$ gives an upper limit and $c_s = \sqrt{kT_d/m_d}$ denotes the dust sound velocity), we obtain for the accretion time $t_{\mathrm{acc}} = D^2/\zeta_{\mathrm{turb}} \approx 5 \times 10^{13}\,\mathrm{s} \gg t_{ff}$. The actual process is expected to work on a time scale somewhere between these two limits.

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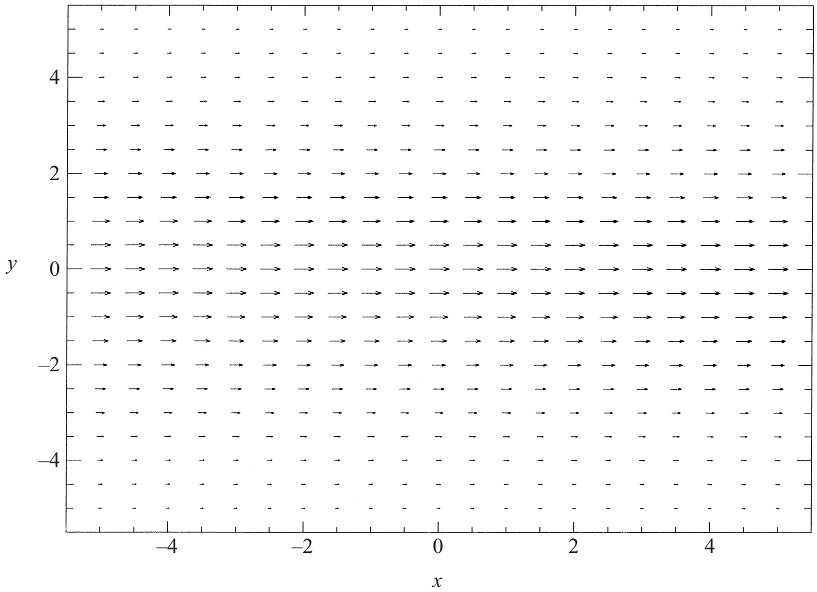

Figure 2. The initially applied dust flow profile.

The fate of the generated magnetic field depends on the interplay of dissipation (due to viscosity and diffusivity) and the battery effect of the ion–dust–neutral system. As mentioned above, the generation time is given by the inverse effective collision frequency ν_{coll}^{-1}, i.e. the field generation at least operates on the time scale of the inverse ion–neutral collision frequency, which is $\nu_{in}^{-1} = 10^{-3}$ s for the parameters chosen. The dissipation in a turbulent accretion disk cannot be faster than $t_{\text{diss}} = L_z^2/\zeta_{\text{turb}} = 10^7$ s for our example. Moreover, the generation has to be faster than the accretion, i.e. $\nu_{\text{coll}}^{-1} < t_{\text{acc}}$. It is clear that, for our applications, these conditions are easily fulfilled, and we can conclude that, notwithstanding the present uncertainties of the actual plasma parameters, our proposed magnetic field generation mechanism is so efficient that circumstellar disks around T Tauri stars should be effectively magnetised after a very short period compared with their lifetimes.

4. Numerical studies

Since the above discussion is restricted to a kinematic treatment – a prescribed velocity field is assumed and merely the influence on the magnetic field evolution is examined – it seems worthwhile to study the proposed magnetic field generation mechanism using dynamical computer simulations. For this purpose, we make use of the 3D Cartesian multifluid code DENISIS (Schröer et al. 1998; Schröer and Kopp 2000) that integrates the mass, momentum, and inner energy balance equations for the different fluids as well as the induction equation. This code uses an explicit integration scheme of second-order accuracy based on a modified Leapfrog algorithm.

In order to qualitatively illustrate the magnetic field generation mechanism, we start from a homogeneous ($\rho_n : \rho_d = 10 : 1$), isothermal, partially ionised dusty plasma to give an idealised representation of a circumstellar disk. At this stage, we do not model the magnetic field self-generation with realistic numbers. That task

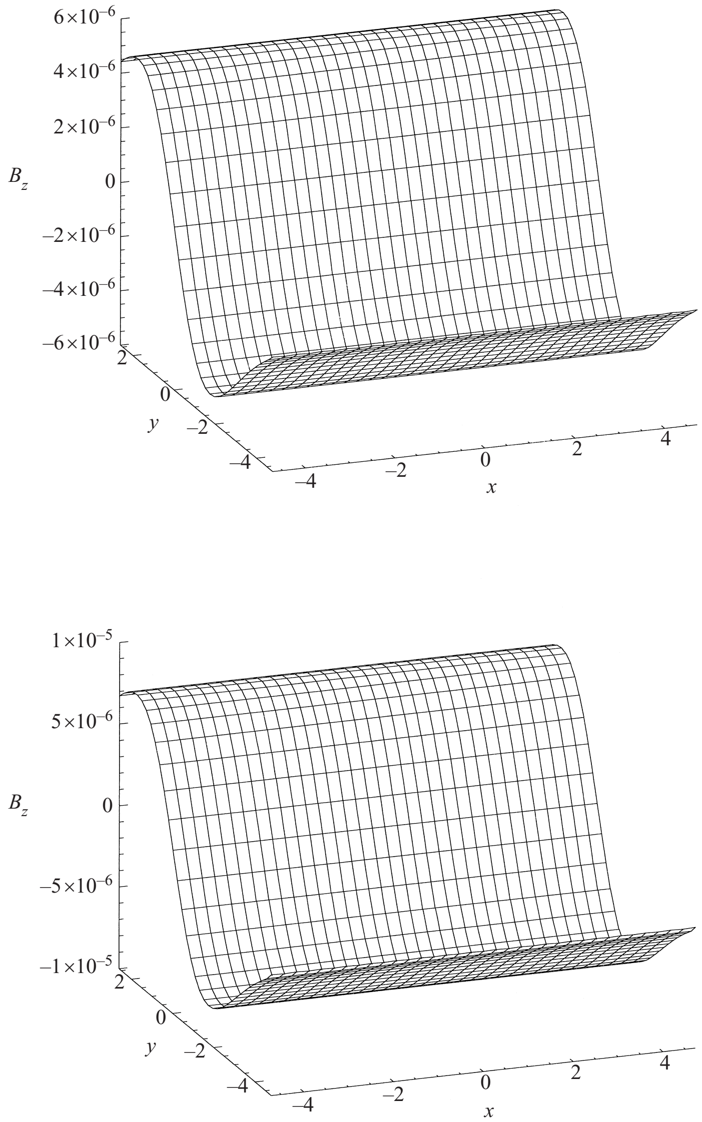

Figure 3. Generation of a magnetic field in the accretion disk after $t = 10$ (a) and $t = 40$ (b) dynamical times. The magnetic field strength is given in units of $B_0 = c_s\sqrt{4\pi\rho_d}$.

demands a detailed numerical simulation study that we plan to carry out in the future. Rather, we intend to show by our numerical study that the field generation mechanism is self-consistently maintained by the relative fluid motion between charged and neutral fluid components.

We choose the frame of the neutral gas component, and thus apply the velocity perturbation for the dust component as shown in Fig. 2. This flow models the relative starward motion of the disk material, and is maintained by the gravitational force. For convenience, we use a Cartesian geometry as a local approximation. The starward accretion flow is directed along the x axis. The disk height extends in the y direction, and the azimuthal direction φ is mapped to the Cartesian z coordinate.

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The relevant collision frequencies are chosen as $\nu_{in} = \nu_{id} = 0.1$, normalised to the inverse characteristic time scale, i.e. the shear length L_y divided by the dust sound velocity c_s. Note that such collision frequencies are very small compared with realistic values for the disk. In Fig. 3, the generated magnetic field is illustrated after (a) 10 and (b) 40 dynamical times L_y/c_s. The simulation results show quasistationary mass transport, i.e. the mass densities remain nearly constant during our evolution time scale. In this numerical treatment, the dynamical back-reaction of the evolving magnetic field on the fluid is taken into account self-consistently, whereas the analytical approach is restricted to a kinematic treatment. It is also worth mentioning that reasonable relative flows will result in magnetic fields associated with electric currents that flow in the circumstellar disks. Since such current sheets are known to be unstable against the resistive tearing mode, this may have implications for the stability and saturation of the disk magnetic field, depending on the local plasma parameters. A detailed numerical study of the dynamics of a circumstellar accretion disk and the associated magnetisation is beyond the scope of this paper, and will be addressed in future work.

5. Conclusions

In partially ionised dusty plasmas, relative flows between dust and neutral components can lead to efficient self-generation of magnetic fields. This generation mechanism has been investigated analytically and numerically with application to accretion disks around T Tauri stars. We can show that these disks are expected to generate their own magnetic field continuously by sheared macroscopic relative dust–neutral motions. This relative radial motion is a consequence of the different coupling properties of the dynamical dust and the neutral component of the accretion disk. Whereas we quantitatively analysed this magnetic field self-generation mechanism for typical T Tauri parameters, it should be noted that magnetisation of the accretion disk should also work efficiently in even earlier phases of stellar evolution (e.g. protostellar class I objects). The necessary condition for this mechanism to work is the existence of a significant stellar magnetic field. This leads to a difference between the corotating and Keplerian motions of the multifluid components that constitute the accretion disks. In particular, we note that the magnetic field generation term does not depend on the sign of the dust charge. The generation process does not saturate due to the existence/generation of magnetic fields, as is the case for the Biermann mechanism (Mestel 1962). As long as the disk material (dust, neutrals, and charged particles) is moving radially, the proposed current and magnetic field generation mechanism will be sustained. In the presence of a small radial magnetic field component, the Balbus–Hawley shear instability (Balbus and Hawley 1991) can result in a faster than linear growth of the poloidal magnetic field component, provided that the appropriate onset criterion is fulfilled. Our findings imply that the interaction between the central stellar objects and the surrounding accretion disks is characterised by rather complex magnetic coupling processes. This magnetic coupling is expected to significantly influence disk accretion (von Rekowski and Fröhlich 1997; Godon 1996), winds and jets (Paatz and Camenzind 1996; Ray et al. 1997), magnetic activity (André 1996; Stelzer and Neuhäuser 2000), and the formation of vortices in circumstellar disks (Adams and Watkins 1995). The magnetic fields in accretion disks may thus prove essential for our understanding of these phenomena in YSOs.

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