NEW COMPOSITE MODELS OF PARTIALLY IONIZED PROTOPLANETARY DISKS

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ABSTRACT

We study an accretion disk in which three different regions can coexist: MHD turbulent regions, dead zones, and gravitationally unstable regions. Although the dead zones are stable, there is some transport due to the Reynolds stress associated with waves emitted from the turbulent layers. We model the transport in each of the different regions by its own $\alpha$ parameter, which is $10^{-3}$ times smaller in dead zones than in active layers. In gravitationally unstable regions, $\alpha$ is determined by the fact that the disk self-adjusts to a state of marginal stability. We construct steady-state models of such disks. We find that for uniform mass flow, the disk has to be more massive, hotter, and thicker at the radii where there is a dead zone. In disks in which the dead zone is very massive, gravitational instabilities are present. Whether such models are realistic or not depends on whether hydrodynamical fluctuations driven by the turbulent layers can penetrate all the way inside the dead zone. If the disk is not in a steady state at some stage of its evolution, then the surface density will evolve toward the steady-state solution. However, if the value of $\alpha$ in the dead zone is much smaller than that in the active zone, the timescale for the parts of the disk that are beyond a few AU to reach a steady state can become longer than the disk lifetime. Steady-state disks with dead zones are a more favorable environment for planet formation than are standard disks, since the dead zone is typically 10 times more massive than a corresponding turbulent zone at the same location.

Subject headings: accretion, accretion disks — planetary systems: protoplanetary disks — stars: pre–main-sequence

1. INTRODUCTION

Protoplanetary disks are believed to have a complex structure. In the parts that are ionized enough (active zones), the magnetic field couples to the gas and the magnetorotational instability (MRI; Balbus & Hawley 1991) develops, leading to turbulence and angular momentum transport. Some parts, however, are too cold and dense for the ionization to reach the level required for MHD turbulence to be sustained. However, in these so-called dead zones, there may still be some low level of transport due to Reynolds or nonturbulent Maxwell stresses (Fleming & Stone 2003; Turner & Sano 2008). Finally, parts of the disk can be dense enough that gravitational instabilities develop, redistributing mass and angular momentum in such a way that the disk settles into a state of marginal stability. To date, there have been no numerical simulations of such global disks. Global three-dimensional MHD simulations so far have modeled stratified and fully turbulent disks, but with no dead zones (Fromang & Nelson 2006).

It is of interest to study whether nonuniform disks composed of different regions as described above can be in a steady state or not. Previous studies (Gammie 1999; Armitage et al. 2001) assumed that no transport at all took place in the dead zone. Therefore, mass could only pile up there until gravitational instabilities developed. If the level of transport is nonzero in the dead zone, however, the disk may be able to adjust to a steady state. It is this possibility that we investigate in this paper.

In determining the characteristics of the disk vertical structure, we ignore reprocessing of the stellar radiation. Only heating associated with the different (magnetic and hydrodynamic) stresses is considered. Reprocessing has been shown to be important when the disk surface is flared (Kenyon & Hartmann 1987; Chiang & Goldreich 1997; D’Alessio et al. 1998; Dullemond et al. 2001). It provides an extra heating source that dominates the disk beyond a few AU and has important consequences for the structure of the dead zone (Matsumura & Pudritz 2006). However, in the models we present below, the parts of the disk beyond 1 AU are shielded from the stellar radiation by the inner parts. These models, which do not include reprocessing, are therefore self-consistent.

The plan of the paper is as follows. In § 2 we describe the way in which we model the transport in a disk with turbulent regions, dead zones, and gravitationally unstable regions. In § 3 we present the equations governing the disk’s vertical structure and describe how they are solved. In § 4 we present the results of the calculations, which we then discuss in § 5. We find that steady-state solutions exist and that they correspond to disks that are thicker, hotter, and more massive than disks without dead zones.

2. MODELING OF THE DISK TRANSPORT

Protoplanetary disks are almost certainly magnetized, as they form from molecular clouds, which have been observed to contain a magnetic field. If the magnetic pressure is smaller than the thermal pressure, then the parts of the disk where the field couples to the gas are prone to the MRI. Its nonlinear development is turbulence that transports angular momentum outward and therefore enables accretion (Hawley et al. 1995; Brandenburg et al. 1995). Recent numerical simulations of the MRI in shearing boxes have shown that the level of transport depends on the Prandtl number (the ratio of the viscosity to the resistivity; Fromang et al. 2007; Lesur & Longaretti 2007). The simulations, which have not converged yet, also show that the turbulence is not sustained when the magnetic flux is zero and for the range of parameters investigated (relatively large Prandtl numbers). It is reasonable to suppose that the magnetic field in protoplanetary disks does not vary very much on the scale of a few scale heights. For this reason, and if these disks were to be described locally by a shearing box, it would be more appropriate to assume a finite net flux. In the parts of the disk that are coupled to the field, we will therefore assume

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that turbulence is sustained and quantify the angular momentum transport using the parameterization of Shakura & Sunyaev (1973), with \( \alpha_T \) being the ratio of the stress to the pressure (the subscript “T” refers to turbulence). It has been shown that this prescription, which is formally valid only if the disk evolves under the action of a viscosity, can be used to describe the global properties and evolution of MHD turbulent disks, although details and stability issues cannot be studied within this framework (Balbus & Papaloizou 1999). Given that \( \alpha_T \) as measured in the simulations depends on the Prandtl number, and that the Prandtl numbers typical of protoplanetary disks are far too small to be investigated numerically, we will assign values to \( \alpha_T \) that are based on observations. The diffusion timescale is indeed limited by the disk lifetime, which is observed to be of a few \( \times 10^9 \) yr. This implies values of \( \alpha_T \sim 10^{-3} \ldots 10^{-2} \) (at least in the outer parts of the disk, which are fully turbulent).

In protoplanetary disks, the minimum ionization fraction that is required for MHD turbulence to be initiated and sustained is on the order of \( 10^{-13} \) on scales of 1 AU (e.g., Balbus & Hawley 2000). Although very small, this level of ionization is not reached in some parts of the disks, which are relatively cold and dense. The extent of the dead zone has been computed by several authors, taking into account ionization by cosmic rays (Gammie 1996; Sano et al. 2000, who also included radioactivity) and X-rays (Igea & Glassgold 1999; Fromang et al. 2002; Matsumura & Pudritz 2003, who also included cosmic rays and radioactivity). In this paper, we will denote by \( \Sigma_T \) the surface density of the active layer; i.e., the surface density beyond which ionizing photons cannot penetrate. Depending on the energy of these photons, \( \Sigma_T \) may vary in the range \( 10^{10} \ldots 10^{11} \) g cm\(^{-2}\) (Gammie 1996; Fromang et al. 2002). Note that the ionization in the disk inner parts is thermal, with the most abundant ions being Na\(^+\) and K\(^+\). For the ionization fraction to be larger than \( 10^{-13} \), the temperature has to be higher than some value \( T_e \), which is close to \( 10^5 \) K over a large range of densities (Balbus & Hawley 2000). There is evidence from numerical simulations that some level of transport can be present in a dead zone that is sandwiched in between turbulent layers (Fleming & Stone 2003), as the turbulence drives hydrodynamic waves that propagate in the dead zone and are associated with a Reynolds stress. More recently, Turner & Sano (2008) have also proposed that magnetic field should be able to diffuse from the active layers into the dead zone. Although the resistivity there is too high for MHD turbulence to be sustained, there is a large-scale magnetic stress associated with the field that drives accretion. (If dust grains are present in the disk, however, the magnetic field may not be able to diffuse down to the disk midplane at all radii.) The transport and energy deposition in the dead zone due to these stresses are nonlocal. However, for illustrative purposes, we will parameterize this transport with a Shakura & Sunyaev (1973) parameter that we will denote by \( \alpha_D \) (where the subscript “D” refers to the dead zone). Fleming & Stone (2003) found that the Reynolds stress in the dead zone is about 10\(\%\) of the Maxwell stress in the active layers. Normalized by the thermal pressure, this translates into \( \alpha_D \) being between 0.07\(\alpha_T \) and 0.3\(\alpha_T \). However, in their simulations, the mass of the dead zone was only a few times that of the active layer. We might expect to find lower values of the Reynolds stress deep into the dead zone when the mass of the active layer is significantly smaller than that of the active zone, which is the case in the models we present below. Therefore, we will take \( \alpha_D \) to be in the range \( 10^{-3} \alpha_T \sim 10^{-4} \alpha_T \).

In the parts of the disk that are dense and cold, gravitational instabilities develop. Spiral density waves grow and transport angular momentum outward and mass inward. The level of saturation is determined by the fact that the disk settles into a state of marginal stability, with the Toomre parameter \( Q \) being on the order of 1.5 (Laughlin & Bodenheimer 1994; Lodato & Rice 2004). It was shown by Balbus & Papaloizou (1999) that in general, in a gravitationally unstable disk, the transport is nonlocal and cannot be described using the viscous disk theory. However, these authors pointed out that a local transport model might apply if the disk maintains itself in a state of marginal stability, which is the case considered here. In the regions where the disk is gravitationally unstable, we will therefore use the Shakura & Sunyaev (1973) prescription and will adjust the transport parameter, which we will denote by \( \alpha_G \) (where the subscript “G” refers to gravity), so as to have \( Q \sim 1.5 \).

It is important to keep in mind that here the disk is modeled in a very schematic way. We believe that this does not affect the gross features of the models we present below, although the details should be considered with caution.

3. DISK VERTICAL STRUCTURE

We consider a system of cylindrical coordinates \((r, \varphi, z)\) with origin at the central star, where \( z = 0 \) is the disk midplane. In our notation, \( P \) is the pressure, \( \rho \) is the mass density per unit volume, \( T \) is the temperature, \( \Omega \) is the angular velocity, \( \nu \) is the “enhanced” kinematic viscosity associated with \( \alpha_T \) in turbulent layers, \( \alpha_D \) in dead zones, and \( \alpha_G \) in gravitationally unstable regions, and \( \kappa \) is the opacity, which in general depends on both \( \rho \) and \( T \). The thin-disk approximation is used throughout (and is checked to be valid a posteriori), so that \( \Omega^2 = GM_*/r^3 \), where \( M_* \) is the mass of the central star and \( G \) is the gravitational constant.

The disk vertical structure is described by the equation of vertical hydrostatic equilibrium,

\[
\frac{1}{\rho} \frac{\partial P}{\partial z} = -\Omega^2 z, \tag{1}
\]

together with the energy equation, which states that the rate of energy removal by radiation is locally balanced by the rate of energy production by viscous dissipation:

\[
\frac{\partial F}{\partial z} = \frac{9}{4} \rho \nu \Omega^2 = \frac{9}{4} \alpha \Omega P, \tag{2}
\]

where \( F \) is the radiative flux of energy through a surface of constant \( z \). In order to write the last term of this equation, we have used the relation \( \nu = \alpha c_s^2/\Omega \), where \( c_s = (P/\rho)^{1/2} \) is the isothermal sound speed and \( \alpha = \alpha_T \) in active layers, \( \alpha = \alpha_D \) in dead zones, and \( \alpha = \alpha_G \) in the gravitationally unstable parts of the disk. Therefore, \( \alpha \) is a function of both \( z \) and \( r \). The flux \( F \) is given by

\[
F = \frac{16\sigma T^3}{3\kappa \rho} \frac{\partial T}{\partial z}, \tag{3}
\]

where \( \sigma \) is the Stefan-Boltzmann constant. In principle, equation (3) is valid only at radii where the disk is optically thick. However, when the disk is optically thin, i.e., when the value of \( \kappa \rho \) integrated over the disk thickness is small compared to unity, the temperature gradient given by equation (3) is small, so the results we get are consistent in that case as well. To close the system of equations, we adopt the equation of state of an ideal gas,

\[
P = \frac{\rho k T}{2m_H}, \tag{4}
\]
where $k$ is the Boltzmann constant and $2m_{H_2}$ is the mass of the hydrogen molecule, which is the main component of protostellar disks at the temperatures and densities of interest here.

### 3.1. Boundary Conditions

We have to solve three first-order ordinary differential equations for the three variables $F$, $P$, (or equivalently $\rho$), and $T$ as a function of $z$ at a given radius $r$. Accordingly, we need three boundary conditions at each value of $r$. These have been described in detail in Papaloizou & Terquem (1999), so here we just recall briefly their expression. We denote with a subscript $\text{B}$ values at the disk surface. The flux at the surface is given by

$$F_s = \frac{3}{8\pi} \dot{M} \Omega^2,$$

where $\dot{M} = 3\pi \langle \nu \rangle \Sigma$, in which $\Sigma = \int_H^0 \rho dz$ is the disk surface mass density and $\langle \nu \rangle = \int_H^0 \rho \nu dz \Sigma/\Sigma$ is the vertically averaged viscosity. If the disk is in a steady state, the value of $\dot{M}$ does not vary with $r$ and is the constant accretion rate through the disk. The quantities $H$ and $-H$ are the upper and lower boundaries of the disk, respectively. The surface pressure is given by

$$P_s = \frac{\Omega^2 H \tau_{ab}}{\kappa_s},$$

where $\tau_{ab}$ is the optical depth above the disk. Since we have defined the disk surface such that the atmosphere above the disk is isothermal, we have to take $\tau_{ab} \ll 1$. Provided that this is satisfied, the results do not depend on the value of $\tau_{ab}$ that we choose (see Papaloizou & Terquem 1999). Finally, the surface temperature satisfies the following equation:

$$2\sigma (T_s^4 - T_h^4) = \frac{9\alpha_s k T_s \Omega}{8\mu m_{H_2} \kappa_s} - \frac{3}{8\pi} \dot{M} \Omega^2 = 0.$$
the surface mass density, the optical thickness through the disk, and the midplane temperature hardly vary with $\tau_{ab}$. This is because the mass is concentrated toward the disk midplane in a layer whose thickness is independent of $\tau_{ab}$.

4. RESULTS

In Figure 1, we plot $H/r$ versus $r$ (from 0.05 to 100 AU) for $\dot{M} = 10^{-8} M_\odot$ yr$^{-1}$, $\alpha_T = 10^{-2}$, and $\alpha_D = 10^{-4}$ (left) and $\alpha_T = 10^{-2}$ and $\alpha_D = 10^{-3}$ (right). The solid and dot-dashed curves correspond to standard disk models (no dead zones) with constant values of $\alpha = \alpha_T$ and $\alpha = \alpha_D$, respectively. The dotted, long-dashed, and short-dashed curves correspond to the disk models with values of $\Sigma_T = 10^2$, 50, and 10 g cm$^{-2}$, respectively. In the regions where there is a dead zone, the disk is thicker, more massive, and hotter.

In this model, there is a range of radii at which the mass of the dead zone is much larger than that of the active layer. We denote with $\Sigma_a$ and $\Sigma_d$ the column densities of the active layer and the dead zone, respectively. At $r = 1$ AU, where the dead zone is vertically very extended, we have $\Sigma_a/\Sigma_d = 0.6 \times 10^{-2}$ and $4 \times 10^{-2}$ for $\Sigma_T = 10$ and 50 g cm$^{-2}$, respectively. Whether hydrodynamic fluctuations driven by the turbulence in the active layers can penetrate down to the disk midplane in that case is not clear.
In order to get a model with higher values of \( \alpha \), we have investigated the case in which \( T = 10^{0.2} \) and \( D = 10^{0.3} \), which is the same as above. This is close to the values found by Fleming & Stone (2003) in their simulations, in which the value of \( \alpha = \alpha_T \) and \( \alpha = \alpha_D \), respectively. The dotted, long-dashed, and short-dashed curves correspond to the disk models with values of \( \Sigma_T = 10^2, 50, \) and \( 10^3 \) cm\(^{-2}\), respectively. In the left plots, the solid and dotted lines are indistinguishable. The vertical structure of the disk depends on the vertical extent of the dead zone.

In Figure 3, we plot \( H/r, \Sigma \), and the midplane temperature \( T_m \) versus \( r \) for the same parameters as in Figures 1 and 2. For comparison, we also plot the curves corresponding to standard disk models (no dead zones) with constant values of \( \alpha = \alpha_T \) and \( \alpha = \alpha_D \). As expected, in the parts where there is a dead zone, the disk is thicker, more massive, and hotter. This is because the vertically averaged value of \( \alpha \) decreases as the vertical extent of the dead zone increases, and \( H/r, \Sigma \), and \( T_m \) increase as \( \alpha \) decreases (e.g., Frank et al. 1992). We see that for low values of \( \alpha_T \), i.e., when the ionizing photons do not penetrate deep into the disk, the turbulent layer is thin and the properties of the disk at the radii where there is a dead zone are very close to those of a standard disk with a constant value of \( \alpha = \alpha_D \). For all the values of \( \alpha_T \), \( H/r \) has a maximum at a fraction of or around 1 AU, which produces a puffing up of the disk. This is because the mass density is very high there, so the dead zone is vertically extended, which makes the disk hotter. Note that a maximum is also present in the standard disk models with constant \( \alpha \). The profiles of \( H/r, \Sigma \), and
Fig. 5.—Same as the bottom panel of Fig. 1, but for \( \alpha_T = 10^{-3} \) and \( \alpha_D = 10^{-5} \). Here the disk is gravitationally unstable from \( r \sim 3 \) AU all the way up to the outer radius.

\( T_m \) tend to be a bit steeper than those in standard disk models with no dead zone, as they are bracketed by the profiles corresponding to models with constant values of \( \alpha = \alpha_T \) and \( \alpha = \alpha_D \).

In Figure 4, we plot \( \rho \) and \( T \) as a function of \( z/r \) at \( r = 1 \) AU for \( \alpha_D = 10^{-4} \) and at \( r = 0.4 \) AU for \( \alpha_D = 10^{-3} \) for the same parameters as in Figure 3. When the dead zone extends down to the disk midplane, the disk structure is very close to that of a standard disk with a constant value of \( \alpha = \alpha_D \) for the lowest value of \( \Sigma_T \). For the largest values of \( \Sigma_T \), the dead zone is small and the disk is similar to a standard disk with a constant value of \( \alpha = \alpha_T \). For \( \alpha_D = 10^{-4} \), we actually cannot distinguish between the two models for \( \Sigma_T = 10^2 \) g cm\(^{-2}\). We see that the mass density in the dead zone is typically 10 times larger than that of a standard fully turbulent disk with a constant value of \( \alpha = \alpha_T \).

In Figure 5, we plot \( H/r \) versus \( r \) for \( M = 10^{-8} M_\odot \) yr\(^{-1} \), \( \alpha_T = 10^{-3} \), \( \alpha_D = 10^{-2} \), and \( \Sigma_T = 50 \) g cm\(^{-2}\); i.e., the same values as above, except for lower values of \( \alpha_T \) and \( \alpha_D \). We checked that decreasing \( \alpha_T \) and \( \alpha_D \) by a factor of 10 has almost exactly the same effect as increasing \( M \) by the same factor, with all the other parameters being kept fixed. When \( \alpha_T \) and \( \alpha_D \) are smaller or \( M \) is larger, the disk is more massive and therefore gravitational instabilities develop in a more extended region, up to the disk’s outer boundary. To keep \( Q \) constant at \( \sim 1.6 \) there, the value of \( \alpha_G \) has to reach \( 2.5 \times 10^{-3} \) in the disk’s outer parts; i.e., \( \alpha_G \) has to become larger than \( \alpha_T \).

We have also run models with \( \alpha_T = 10^{-2} \) and \( \alpha_D = 10^{-3} \alpha_T \). We get similar results in that case, with the locations where there is a dead zone being hotter, thicker, and more massive.

We note that in all the models presented above, the puffing up of the inner parts of the disk casts a shadow over the outer parts, which therefore do not reprocess stellar radiation. The models presented here are self-consistent in the sense that if the disk is not illuminated by the star to begin with, it will not be able to reprocess the stellar radiation at any time. There may, however, exist another type of solution in which stellar irradiation is important. Indeed, if the outer parts of the disk are illuminated at some point (before a steady state is reached, for instance), a flaring is produced that could be maintained. To test this hypothesis, we have modified the surface temperature. Instead of using equation (7) to calculate \( T_T \), we fix \( T_T = 300 (r/\text{AU})^{-1/4} \) K. This corresponds to \( T_T \sim 950 \), 300, and \( \sim 95 \) K at \( r = 0.1 \), 1, and 10 AU, respectively. These values are close to the surface temperatures computed by D’Alessio et al. (1999) in irradiated disk models for \( M = 10^{-8} M_\odot \) yr\(^{-1} \) and \( \alpha = 10^{-3} \), this value of \( \alpha \) being intermediate between \( \alpha_T \) and \( \alpha_D \). With this new value of \( T_T \), the profile of \( H/r \) that we get still shows a shadowing of the outer parts. This suggests that even if the disk outer parts are illuminated by the central star at some point before a steady state is reached, they cannot flare enough to remain irradiated after a steady state is established, and so reprocessing of stellar radiation does not play a role in the steady-state models.

5. DISCUSSION AND CONCLUSION

In this paper, we have considered a disk in which transport is produced by either (1) Maxwell stress resulting from MHD turbulence in the (gravitationally stable) regions where the gas is well coupled to the magnetic field, (2) gravitational stress in the parts that are gravitationally unstable, or (3) Reynolds stress associated with hydromagnetic waves driven by the turbulence in adjacent layers and Maxwell stress associated with large-scale field siphoned off from these turbulent layers in the parts that are not turbulent and are gravitationally stable (dead zones). We have modeled the transport using the Shakura & Sunyaev (1973) prescription, with \( \alpha_D \) being \( 10^{-3} \) times smaller than \( \alpha_T \), and with \( \alpha_G \) being adjusted so as to give a Toomre parameter of \( Q \sim 1.5 \) in the gravitationally unstable regions. Although such a modeling is in principle not valid when the transport is nonlocal, we believe that it does not invalidate the gross features of the models we have presented.

We have found that steady-state models of such a disk exist, and that they are physically reasonable. Since \( M = 3\pi \nu \Sigma \alpha c \Sigma \) mass flow through the disk can be uniform only if the disk is more massive, hotter, and thicker at the radii where the vertically averaged value of \( \alpha \) is smaller; i.e., where there is a dead zone. Note that even though the disk is thicker at the locations where there is a dead zone, it remains thin for the parameters investigated here. In disks in which the dead zone tends to be massive, which is the case for the lowest values of \( \alpha_D \) and/or the lowest values of \( \Sigma_T \) investigated here, gravitational instabilities control the dynamics of part of the disk.

Whether these models are realistic or not depends on whether hydrodynamical fluctuations driven by the turbulent layers can penetrate all the way inside the dead zone. This may be more easily achieved when the ratio of the mass of the dead zone to
that of the active layer is the smallest, which in our models corresponds to a value of \( \alpha_p / \alpha_T = 0.1 \).

If the disk is, at some stage of its evolution, out of a steady state, then the surface density will change in such a way as for the disk to evolve toward a steady state. To see this, let us consider a disk in which \( M \) decreases inward at some location. Then mass will accumulate there, as accretion through this region is slower, until \( \Sigma \) is large enough for the flow to be steady. In contrast, if \( M \) increases inward at some location, accretion there will be faster, and mass will be depleted until \( \Sigma \) is reduced to the level where a steady state is reached. The timescale \( t_r \) for establishing a steady state at the radii where there is a dead zone may be long, however, as it is given by

\[
t_r = \frac{r^2}{3(\rho)} \sim \frac{1}{3(\alpha)} \left( \frac{r}{H} \right)^2 \Omega^{-1},
\]

where \( \alpha = \int_H^\infty \rho \, dz / \Sigma \). With values of \( \alpha = 10^{-2} \) and \( H/r = 0.1 \), we find that \( t_r \sim 5 \times 10^7 \) and \( 6 \times 10^8 \) yr at 1 and 5 AU, respectively, which is much smaller than the disk lifetime. However, if the dead zone is vertically very extended at these radii, we have \( \alpha \sim \alpha_p \), so that the viscous timescale can get 100 times longer if \( \alpha_p = 10^{-4} \). At 1 AU we still get a timescale that is significantly smaller than the disk lifetime, but this is not the case at 5 AU, where the disk may not be able to reach a steady state. If these parts of the disk build up from mass that flows in from farther out, then the mass in the dead zone will slowly increase. This process will never produce outbursts as envisioned by Gammie (1999), though. Indeed, enough mass could pile up for an outburst to be produced, a steady state would be achieved in which mass would be transported either by density waves or Reynolds stresses driven by the turbulent layers.

Note that the spectral energy distribution of a steady disk such as those considered in this paper is not different from that of a standard disk with no dead zone, as the total flux emitted at the surface depends only on \( M \). The surface temperature also does not depend on \( \Sigma_r \), as can be seen from equation (7). However, the temperature just below the surface is significantly larger in disks with dead zones than in standard disks. The chemistry there would therefore be different. Grain properties may also differ. To study this, irradiation of the disk by the central star would have to be taken into account. Although it will not be important beyond the regions where there is a dead zone, as these parts are shielded from the stellar radiation due to the puffing up of the disk, it will be important in the parts where \( H/r \) increases.

Finally, we comment that a steady-state disk with a dead zone is a more favorable environment for planet formation than is a standard disk, as the dead zone, which in general encompasses the region of planet formation, is significantly more massive. Planet formation would therefore be faster there (Lissauer 1993). Note, however, that there is an issue as to how to stop type I migration of the cores in these dead zones. Turbulence, which has been shown to alter this type of migration (Nelson & Papaloizou 2004), cannot be invoked. But if large-scale fields are siphoned off from the turbulent layers in these regions, they may prevent the cores from migrating too fast (Terquem 2003; Fromang et al. 2005; Muto et al. 2008). Planetary cores may also be stopped at the interface between the dead zone and the inner turbulent region, where the surface mass density varies sharply, as suggested by Masset et al. (2006). This process, however, relies on the (positive) corotation torque becoming dominant over the Lindblad torque due to the density gradient and was studied in the context of a laminar viscous disk. It is not clear how the corotation torque is affected by the turbulence at the transition between the active and dead zones and whether it would still be as efficient in this context.

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