GLOBAL HYDROMAGNETIC SIMULATIONS OF A PLANET EMBEDDED IN A DEAD ZONE: GAP OPENING, GAS ACCRETION, AND FORMATION OF A PROTOPLANETARY JET

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ABSTRACT

We present global hydrodynamic (HD) and magnetohydrodynamic (MHD) simulations with mesh refinement of accreting planets embedded in protoplanetary disks (PPDs). The magnetized disk includes Ohmic resistivity that depends on the overlying mass column, leading to turbulent surface layers and a dead zone near the midplane. The main results are: (1) the accretion flow in the Hill sphere is intrinsically three-dimensional for HD and MHD models. Net inflow toward the planet is dominated by high-latitude flows. A circumplanetary disk (CPD) forms. Its midplane flows outward in a pattern whose details differ between models. (2) The opening of a gap magnetically couples and ignites the dead zone near the planet, leading to stochastic accretion, a quasi-turbulent flow in the Hill sphere, and a CPD whose structure displays high levels of variability. (3) Advection of magnetized gas onto the rotating CPD generates helical fields that launch magnetocentrifugally driven outflows. During one specific epoch, a highly collimated, one-sided jet is observed. (4) The CPD’s surface density is $\sim 30$ g cm$^{-2}$, small enough for significant ionization and turbulence to develop. (5) The accretion rate onto the planet in the MHD simulation reaches a steady value $8 \times 10^{-3}$ $M_J$ yr$^{-1}$ and is similar in the viscous HD runs. Our results suggest that gas accretion onto a forming giant planet within a magnetized PPD with a dead zone allows rapid growth from Saturnian to Jovian masses. As well as being relevant for giant planet formation, these results have important implications for the formation of regular satellites around gas giant planets.

\textit{Key words:} magnetohydrodynamics (MHD) – methods: numerical – planets and satellites: formation – protoplanetary disks

\textit{Online-only material:} color figures

1. INTRODUCTION

Gas giant planets are widely believed to form via core nucleated accretion, a scenario that begins with the formation of a solid rock and ice core in a protoplanetary disk (PPD) by the agglomeration of smaller bodies (planetesimals) and which concludes with the accretion of a gaseous envelope from the surrounding nebula (Mizuno 1980). Although it has been suggested that giant planets may form through the direct gravitational fragmentation of a massive PPD during its early evolution (Boss 1998), circumstantial evidence for core accretion having operated in the solar system is provided by the inferred existence of significant cores in Saturn, Uranus, (e.g., Saumon & Guillot 2004) and Neptune. It can further be argued that the substantial numbers of relatively low-mass super-Earth and Neptune-like extrasolar planets being discovered (e.g., Lissauer et al. 2011) indicate that core accretion is a common mode of planetary formation outside of the solar system.

Detailed one-dimensional (1D) models of gas giant planet formation indicate that envelope accretion occurs in two distinct stages: (1) a quasi-static contraction phase during which the envelope mass grows slowly over timescales $\gtrsim 1$ Myr (Pollack et al. 1996) and (2) a runaway growth phase during which the envelope accretes dynamically onto the planet. This later phase normally arises once the envelope exceeds the core mass, corresponding to a total planet mass $\gtrsim 35 M_J$. Prior to runaway gas accretion, the protoplanet remains embedded in the nebula and the bloated envelope is envisaged to connect smoothly onto the surrounding disk. During the runaway phase, however, the planet contracts down to a size $\lesssim 3$ Jupiter radii and gas accretion is expected to occur through a circumplanetary disk (CPD) that forms by the flow of gas into the planet Hill sphere (Papaloizou & Nelson 2005). We note that the 1D envelope calculations indicate that planets with Saturn-like masses ($\sim 0.2–0.6 M_J$) are likely to experience maximal gas accretion (Pollack et al. 1996), as the flow onto the planet is relatively unimpeded by the compressional heating of the envelope for these masses. Although these 1D quasi-static calculations cannot determine the details of the hydrodynamic (HD) flow, so some doubt remains about the actual accretion rate, it is noteworthy that extrasolar planets with masses similar to Saturn are common,\footnote{Out of 720 confirmed extrasolar planets listed on exoplanets.org, approximately 170 have masses in the range 0.2–1 $M_J$.} even though existing calculations suggest that once a planet enters the runaway gas-accretion phase it should grow rapidly beyond this mass if it is in the presence of a significant gas reservoir. One motivation for performing the calculations presented in this paper is to address this issue and to examine whether or not a possible bottleneck exists that can prevent rapid growth of planets in this mass range.

The formation of a gas giant planet through gas accretion and the action of tidal torques lead to the formation of an annular gap around the vicinity of the planet (Lin & Papaloizou 1986; Bryden et al. 1999; Kley 1999; Lubow et al. 1999). This may also herald the transition of the planet’s migration from type I (Ward 1997) during the embedded phase to type II when a gap has formed (Lin & Papaloizou 1986; Nelson et al. 2000). Material then feeds onto the planet through the gap at the viscous gap has formed (Lin & Papaloizou 1986; Nelson et al. 2000). Material then feeds onto the planet through the gap at the viscous gap has formed (Lin & Papaloizou 1986; Nelson et al. 2000). Material then feeds onto the planet through the gap at the viscous
supply rate. A detailed understanding of this accretion process necessitates a sophisticated model of the PPD environment in which the planet is embedded.

Global, multidimensional studies of gas accretion onto embedded planets began with two-dimensional flat, viscous, locally isothermal disk models (Bryden et al. 1999; Kley 1999; Lubow et al. 1999) and suggested that accretion rates onto Jovian mass planets should be $\sim 10^{-5} M_J \text{yr}^{-1}$ for disk models with masses close to the minimum mass solar nebula (Hayashi 1981) and canonical values for the disk aspect ratio ($H/r \sim 0.05$) and viscous stress parameter ($\alpha \sim 10^{-3}$). Models that adopt more realistic equations of state and/or are three-dimensional (3D) find similar accretion rates (D’Angelo et al. 2003; Klahr & Kley 2006; Paardekooper & Mellema 2008; Ayliffe & Bate 2009b).

Typically these simulations do not resolve the flow in the planet Hill sphere very accurately, so although they indicate the presence of a CPD, the simulations are unable to resolve the details of its flow. An exception to this are the calculations by Ayliffe & Bate (2009b), which attempt to simulate the full 3D radiation-hydrodynamic evolution of the planetary envelope and CPD, but unfortunately these are hampered by time-step restrictions that prevent long-term runs from being performed.

Recent 3D HD studies of the accretion flow onto giant planets have been performed that focus on the local evolution in the vicinity of the Hill sphere (Machida et al. 2008; Taniwaga et al. 2012). These highly resolved studies have uncovered a particularly interesting result, namely that accretion onto the planet occurs not through the midplane region of the CPD—where the gas flow appears to be away from the planet on average—but instead through gas flows that occur at higher latitudes within the planet Hill sphere. The midplane region of the CPD does not appear to act as a traditional accretion disk, but instead is a region where material that enters the Hill sphere with excess angular momentum is spun-out away from the planet. The most recent study of Taniwaga et al. (2012) resolves the flow down to approximately 3% of the Hill sphere radius and midplane outflow appears to occur all the way down to this region. It is unclear how these results will change when angular momentum transport processes are included in the evolution of the CPD.

All of the above multidimensional studies have either ignored the disk viscosity or have solved the Navier–Stokes equations by adopting an anomalous value to account for the angular momentum transport arising from turbulence in the disk, which is believed to be driven by the magnetorotational instability (MRI; Balbus & Hawley 1998). Highly resolved studies of gas accretion onto giant planets embedded in disks that support MRI turbulence have not been performed, although low-resolution studies of gap formation and/or gas accretion have been presented by Nelson & Papaloizou (2003), Winters et al. (2003), Papaloizou et al. (2004), and Uribe et al. (2011) for disks that support fully developed turbulence. It was observed that magnetic braking caused an apparent increase in gas accretion in some of these simulations by removal of angular momentum of material entering the Hill sphere, but the low resolution and assumption of ideal magnetohydrodynamics (MHD) render this observation questionable in terms of its application to real systems. Although these simulations in some ways represent a step-up in realism relative to laminar viscous disk models, the cold, dense midplane regions of PPDs are believed to host dead zones where the magnetic field and gas decouple due to the low ionization levels there, with accretion occurring in the surface layers that are ionized by external sources such as cosmic rays (CRs) and stellar X-rays (XRs; Gammie 1996; Igea & Glassgold 1999).

In this paper, we present the first global simulation of a giant planet embedded in a magnetized PPD with a midplane dead zone and actively accreting surface layers, where the larger-scale flow features in the planet Hill sphere are resolved. We adopt a resistivity prescription in which the ionization structure changes over time in response to the changing density, temperature, and column of material absorbing the XRs (Gressel et al. 2011, 2012) and examine gap formation and accretion onto the protoplanet. We utilize adaptive mesh refinement to provide good resolution within the planet Hill sphere, resolving the flow with reasonable accuracy down to a distance from the planet equal to 5% of the Hill sphere radius. Similarly resolved HD simulations are also presented for comparison purposes. The primary aim is to examine the structure of the gap that forms as the planet accretes gas and exerts tidal torques on the surrounding PPD (including the influence of the varying ionization fraction as the gap opens) and to examine the flow in the Hill sphere and the rate of gas accretion onto the planet. We find accretion rates that are consistent with the previous studies described above, implying that runaway gas accretion should form planets with Jovian masses and above efficiently if planet formation occurs in a disk with mass similar to the minimum mass nebula. Enlivening of the gap region into a turbulent state leads to a highly time-dependent 3D flow within the gap and planet Hill sphere, with interesting consequences for the dynamics of the CPD.

This paper is organized as follows: in Section 2, we formulate the equations and lay out the numerical methods used. The utilized disk model is described in Section 3, where we give initial and boundary conditions (BCs) and recapitulate our ionization model. General results concerning the PPD are presented in Section 4, followed by Section 5 on the opening of the gap and Section 6 on the CPD. Results on the accretion flow onto the planet are found in Section 7. We finally summarize our findings in Section 8, where we discuss potential implications of our results.

2. NUMERICAL METHODS

We perform HD and MHD simulations of protoplanetary accretion disks employing a spherical-polar mesh with adaptive grid refinement around an embedded planetary core. The planet is modeled via the gravitational potential of a softened point mass and its position is kept fixed throughout each simulation. The planet’s mass is allowed to grow during the simulation by accretion of gas from the disk, as detailed below.

2.1. Numerical Scheme

The simulations presented in this paper were performed using the single-fluid MHD code NIRVANA-III, which is based on a second-order, finite-volume Godunov scheme (Ziegler 2004) and employs the constrained transport (CT) discretization for an intrinsically divergence-free evolution of the magnetic induction. The code has recently been extended to orthogonal curvilinear meshes (Ziegler 2011) and here we use spherical-polar coordinates ($r, \theta, \phi$), denoting spherical radius, co-latitude, and azimuth, respectively. Deviating from the publicly available version of the code, we here use the upwind reconstruction technique of Gardiner & Stone (2008) to obtain the edge-centered electromotive force (EMF) needed within the CT update. We have furthermore generalized the EMF interpolation to
curvilinear coordinates (see Skinner & Ostriker 2010, for the cylindrical case). The upwind reconstruction avoids stability issues present in the original EMF interpolation scheme by Balsara & Spicer (1999) and the relevance of this to the development of the MRI was demonstrated by Flock et al. (2010). By default, Nirvana-iii addresses this issue by implementing the two-dimensional Riemann solver of Londrillo & del Zanna (2004), but we here instead chose the approach taken by Gardiner & Stone, as this allows us to easily use the more accurate HLLD approximate Riemann solver of Miyoshi & Kusano (2005), which provides better numerical accuracy for low Mach number flows. The benefits of using HLLD when modeling MRI turbulence have been demonstrated by Balsara & Meyer (2010).

2.2. Equations Solved

For our HD calculations, we solve the compressible Navier–Stokes equations subject to a prescribed enhanced viscosity, which we estimate from the turbulent stresses occurring in the MRI simulation. The hydrodynamic run employs the standard resistive MHD equations with a position-dependent molecular diffusivity $\eta(\mathbf{r}, t)$, which we derive self-consistently from a detailed ionization model (see Section 3.3) accounting for irradiation of the disk surface by ionizing sources. The full set of equations in a coordinate system corotating with the planet reads:

$$
\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0,
$$
$$
\frac{\partial \rho \mathbf{v}}{\partial t} + \nabla \cdot (\rho \mathbf{v} \otimes \mathbf{v} + \rho \mathbf{a} + \nabla \cdot \mathbf{a}) = -\rho \nabla \Phi + \rho \mathbf{a} + \nabla \cdot \mathbf{F},
$$
$$
\frac{\partial \rho e}{\partial t} + \nabla \cdot (\rho e \mathbf{v} - (\rho \mathbf{B} \otimes \mathbf{B}) \mathbf{v} - \rho \mathbf{a} \cdot \mathbf{v}) = -\rho \nabla \Phi \cdot \mathbf{v} + \rho \mathbf{a} \cdot \mathbf{v} + \nabla \cdot \left(\mathbf{v} \otimes \rho \mathbf{B} \otimes \mathbf{B} - \mathbf{J} \right) + \Gamma,
$$
$$
\frac{\partial \mathbf{B}}{\partial t} - \nabla \times \left(\mathbf{v} \otimes \mathbf{B} - \eta \mathbf{J}\right) = 0,
$$
$$
\nabla \cdot \mathbf{B} = 0,
$$

with conserved variables $\rho$, $\rho \mathbf{v}$, and the total energy $e = e + (1/2) \rho \mathbf{v}^2 + (1/2) \mathbf{B}^2$, and where $e$ denotes the thermal energy density. We use $\mathbf{J} = \text{curl} \mathbf{B}$ and introduce the total pressure $p^* = p + (1/2) \mathbf{B}^2$. The inertial acceleration due to Coriolis and centrifugal effects is $\mathbf{a}_c = -2 \Omega \mathbf{z} \times \mathbf{v} - \Omega \times (\Omega \times \mathbf{r})$, where $\Omega = \Omega_\ast \mathbf{z}$ is the axial vector representing the angular frequency at the planet radius and $\mathbf{z}$ is the unit vector pointing along the rotation axis of the PPD. The gravitational potential is composed of the central potential of a solar-mass star ($\Phi = M_\odot$) and the softened point-mass potential of the planet with mass $M_p = M_p(t)$ at the fixed position $\mathbf{r}_p$, i.e.,

$$
\Phi(\mathbf{r}, t) = -\frac{GM_\odot}{r} - \frac{GM_p}{|\mathbf{r} - \mathbf{r}_p|} + \epsilon \left(\mathbf{r} - \mathbf{r}_p\right)^3 r_p^3.
$$

The third term appearing in the potential is often referred to as the “indirect term” and it accounts for the fact that our coordinate system has its origin at the position of the star, rather than at the center of mass of the combined system. Note that the gravity of the disk does not act on any of the bodies and hence it does not appear in the potential function. Even with refined meshes, we do not yet resolve the physical radius of the accretion envelope of the planet core. To avoid the singularity of the assumed point-mass potential, the second term in Equation (2) is softened with a smoothing length that is fixed at 5% of the Hill radius corresponding to the initial planet mass of $100 M_\oplus$. This corresponds to half a grid spacing on the coarse mesh.

For the cases where we evolve an energy equation, we assume a relation, $p = (\gamma - 1) e$, with $\gamma = 7/5$. The additional heating and cooling term, $\Gamma$, serves the purpose of relaxing the thermal energy $\epsilon$ toward the initial radial temperature profile $T_i(R)$ on a local dynamical timescale. The update is implemented operator-split and applies Newtonian cooling according to the equation

$$
\frac{1}{\rho} \frac{d\epsilon}{dt} = -\frac{\Omega_\ast(R)}{2\pi} \left(\epsilon - \Theta_\ast(R)\right),
$$

where $\Omega_\ast(R)$ is the Keplerian rotation frequency as a function of cylindrical radius, $R$, and the ideal gas relation is used to compute $\Theta_\ast(R) = T_i(R) k_b / (\mu m_\ast (\gamma - 1))$. While such a treatment is a far cry from a realistic approximation of radiative effects occurring in a real PPD, it adds realism over a purely isothermal approach. On a practical level, it successfully suppresses the occurrence of the vertical-shear instability (Nelson et al. 2013), which leads to severe disturbance of the dead-zone layer via the excitation of corrugation waves. Such disturbances were encountered in inviscid simulations of a locally isothermal disk with imposed radial temperature profile, but in the locally isothermal run that we present here the inclusion of viscosity in the model also suppresses the vertical-shear instability.

The viscous stress tensor, $\tau$, appearing in the momentum equation, is given by $\tau = \nu (\nabla \mathbf{v} + (\mathbf{v} \nabla) - (2/3) \nabla \mathbf{v})$, with $\nu$ being the dynamic viscosity parameter. Ohmic dissipation enters the induction equation via the $\nabla \times (\eta \mathbf{J})$ term. In our explicit time-integration scheme, the numerically allowed time step related to dissipative terms scales with the square of the grid spacing. This quickly becomes restrictive when applying mesh refinement, especially in the presence of high values of the viscosity coefficient $\nu$ and the diffusivity $\eta$. To circumvent these potentially very restrictive time-step constraints, we have adopted the super-time-stepping (STS) scheme, introduced by Alexiades et al. (1996). This concept is based on the idea of applying a sequence of (first-order accurate) forward-Euler sub-steps for the dissipation terms. With an appropriate choice of non-equidistant sub-intervals in time, the scheme gains an advantage over classical sub-stepping techniques. We here chose the STS parameter $\nu_{STS} = 0.02$ and limit the maximum ratio between the permissive Courant time step and the STS time step to 20. We remark that the energy equation contains a divergence of a dissipative flux, $\nabla \mathbf{v} \cdot \eta \mathbf{B} \times \mathbf{J}$. Comparing with test solutions obtained with the unmodified Nirvana-iii, we have not found any indications that the combination of STS with an associated forward-Euler update of the energy equation leads to spurious results. We remark that a future implementation of the Nirvana-iii code will make use of the more accurate second-order scheme recently proposed by Meyer et al. (2012).

3. MODEL DESCRIPTION

Our underlying protostellar disk model aims to resemble the intermediate regions of the early protosolar nebula where giant planet formation is believed to have taken place. The adopted computational domain spans a radial extent of $r \in [1, 8]$ au. Owing to limited computational resources, we restrict the azimuthal domain to a quarter wedge, i.e., $\phi \in [0, \pi/2]$. The latitudinal grid spans a region $\theta \in [\pi/2 - \bar{\theta}, \pi/2 + \bar{\theta}]$, with $\bar{\theta} = 4.5 H/r$, i.e., covering four and a half pressure scale heights, $H$, on each side of the disk midplane. Unless otherwise indicated, the base grid resolution is chosen to be $N_r \times N_\phi \times N_\theta = 384 \times 96 \times 128$ grid points, corresponding to 10.7 grid points per $H$ in the vertical direction. This coincides

6 Note that for clarity, we have suppressed factors of the permeability $\mu$ appearing in the equations.
with the requirement to reasonably resolve unstable MRI modes (see the Appendix). Block-adaptive grid refinement (with blocks of size $4 \times 4 \times 4$) is enabled once the planet potential is switched on. Refined meshes of level $l = 1, 2, \ldots$ have grid spacings of $2^{-l}$ times the base-level resolution. For reasons of simplicity, the adaptivity is controlled by a purely geometrical criterion. Approximately spherical regions with $\Delta r_l \propto r$ around the planet are refined with $l = 3, 2, 1$ levels, respectively. This implies that the Hill sphere ($r_c = r_{pl}$) is always refined at the highest level of $l_{\max} = 3$, resulting in a grid resolution of roughly 5 Jovian radii, $R_J$ in the $r$ and $\theta$ directions, and about $11 R_J$ in the $\phi$ direction. Considering the number of grid cells covering the CPD, we note that a planet of mass $M_p = 150 M_\oplus$ located at $r_p = 3.5$ au will be surrounded by a CTD with radius approximately equal to 40% of the Hill radius, giving $r_{CTD} \sim 0.074$ au. The grid spacings on the most refined level are $(\Delta r, \Delta \theta, \Delta \phi) = (2.28, 2.05, 5.37) \times 10^{-3}$ au. In the horizontal plane, $r_{CTD}$ is resolved by approximately 33 cells of length $\Delta r$, and 14 cells of length $\Delta \phi$. Assuming a constant aspect ratio of $H/r \sim 0.5$ of the CTD, the scale height of the CTD is resolved by approximately 18 cells at its outer rim. Toward the center, the resolution per scale height drops linearly. This suggests that the vertical structure of the CPD should be decently resolved, as should the circulating flow in the horizontal plane. Although magnetic stresses related to larger scale magnetic features in the CTD will be resolved, we note that development of the MRI on small scales in the CTD cannot be followed in our simulations.

It is well known that the spherical-polar mesh is ideally suited to numerically preserve angular momentum within the PPD. However, in the vicinity of the planet, the grid is essentially Cartesian—in consequence, the angular momentum of the CTD is likely less well conserved. The associated numerical truncation error will contribute to the measured accretion rate onto the planet.

### 3.1. Equilibrium Disk Model

We now describe the initial disk model, which is chosen to facilitate comparison of our results with the existing literature. Although modeling and observations point to the fact that PPDs are moderately flaring (Watson et al. 2007; Pinte et al. 2008), this results in thin disks at small radii and is suboptimal for simulating using a spherical-polar grid. A convenient setting is obtained when one prescribes a locally isothermal temperature, $T$, falling off with inverse cylindrical radius, i.e., $T(r) = T_0 (r/R_0)^{-1}$. It can easily be shown that such a dependence leads to a constant opening ratio $H/r$ throughout the disk, promoting the use of a spherical-polar domain. If we further prescribe a power-law dependence for the disk midplane density, $\rho_{\text{mid}}(R) = \rho_0 (R/R_0)^{-3/2}$ and assume independent hydrostatic balance in the vertical and radial direction, respectively, we can solve for an equilibrium initial solution given by

$$
\rho(r) = \rho_0 \left( \frac{R}{R_0} \right)^{-3/2} \exp \left( \frac{GM_* c_s^2}{2r} \right), \quad (4)
$$

$$
\Omega(r) = \Omega_K(R) \sqrt{\frac{H}{r}} \frac{5}{2} \left( \frac{H}{R} \right)^2, \quad (5)
$$

where we have introduced the isothermal sound speed $c_s$, which can be derived as $c_s^2 = c_{0s}^2 (R/R_0)^{-1}$. We fix the free parameter $c_{0s}$ according to a value of $h \equiv H/R = 0.05$, noting that $H \equiv c_s \Omega_K$, where the Keplerian angular velocity $\Omega_K(R) = \sqrt{GM_*/R^{3/2}}$. Note that, via $r$, the equilibrium rotation profile has a weak vertical shear. One can also see that in the midplane (i.e., for $r = R$), the flow is sub-Keplerian to order $(H/R)^2$, which is a consequence of radial pressure support. We chose the normalization $\rho_0$ to yield a vertically integrated column density of $\Sigma = 150 \, \text{g cm}^{-2}$ at the location $R = 5$ au, yielding a total disk mass of $M_{\text{disk}} \approx 3.6 M_\oplus$, i.e., when accounting for the full azimuthal extent of $2\pi$. In terms of the disk temperature, we obtain $T = 540 K$ at $R = 1$ au and $T = 108 K$ at a distance of 5 au, comfortably outside of the ice line where water ice condenses. The disk is initially threaded by a weak uniform vertical magnetic field of $B_z = 3.6$ mG. This corresponds to a midplane $\beta_p \equiv 2 p/B^2 \approx 3 \times 10^5$ at the location of the planet.

### 3.2. Boundary Conditions

To facilitate the long-term evolution of our disk model, we have to make modifications to the standard BCs implemented in NIRVANA-III. With the exception of the azimuthal direction, where we simply apply periodicity, all BCs are based on the “outflow” type, i.e., allowing material to leave the domain, but preventing inflow. To improve robustness at the inner corners of the $(r, \theta)$ domain, we impose reflecting boundaries in the $r$ direction beyond 4 $H$ in $\theta$ and moreover impose Keplerian rotation for $v_\phi$.

The BCs applied at the upper and lower disk surfaces for the magnetic field are of the “perfect conductor” type, i.e., enforcing the normal component to be zero at the boundary and applying zero-gradient extrapolation to the parallel field components. In addition, the BCs have been modified to preserve the constant vertical field imposed initially. Because of the vertical stratification, the $\theta$ boundaries require special attention. It is well known that unsplitted finite-volume schemes make it hard to exactly preserve a given static equilibrium (see, e.g., Zingale et al. 2002). This becomes pronounced at the boundaries because using constant extrapolation of the thermal energy density underestimates the pressure gradient. Not balancing gravity exactly will induce a standing accretion shock in the first grid cell of the active domain. To obtain a better estimate for the vertical continuation of the disk profiles outside the active domain, we employ a second-order Runge–Kutta shooting method to integrate the equations of hydrostatic equilibrium, thereby significantly reducing the amplitude of spurious boundary effects.

A similar but somewhat less dramatic effect is present in the radial direction. Taking into account the additional centrifugal force, we extrapolate the radial equilibrium condition to obtain improved boundary values, aiding the long-term stability of the disk evolution. Furthermore, to avoid reflection of spiral density waves, we have implemented “buffer zones” as described in detail in de Val-Borro et al. (2006). This implies that, in narrow annuli adjacent to the radial domain boundary, we force the density $\rho$, energy density $\epsilon$, and the velocity $\mathbf{v}$ back to their initial values. The timescale of this process is proportional to the local dynamical timescale and is smoothly tapered off with increasing separation from the inner and outer radial boundaries.

### 3.3. The Ionization Model

The central aim of this work is to develop realistic MHD models for protoplanetary accretion disks containing gaps opened by embedded planets. This requires that we obtain a self-consistent distribution of the disk’s ionization fraction. For
our new global models, we largely follow the approach taken in the local box models described in detail in Gressel et al. (2011, 2012)—with the exception that the dependence on the radial coordinate is now included explicitly, rather than being evaluated at the reference radius of the box model.

The disk’s ionization balance is strongly affected by the presence of small dust grains (Sano et al. 2000; Ilgner & Nelson 2006). The associated short adsorption timescale of free electrons makes it prohibitive to follow the detailed nonequilibrium chemistry. Instead, we adopt a simplified approach employing a precomputed table based on the reaction network used for model 14 of Ilgner & Nelson (2006). As before, we consider the contributions of all the charged species following Equations (21)–(31) in Wardle (2007). When deriving the resistivity \( \eta(\mathbf{r}, t) \), we restrict ourselves to the case of Ohmic resistivity, which is a reasonable assumption for the denser parts of the PPD (see, e.g., Wardle 2007) under weak magnetic fields. Our choice is mainly motivated by reasons of tractability as the effect of ambipolar diffusion (and Hall EMFs) will greatly increase the computational complexity of the model. Recent work points to the potential importance of including these additional non-ideal MHD effects (Wardle & Salmeron 2012; Bai & Stone 2013) and our longer term intention is to explore these on a case-by-case basis rather than presenting a “kitchen-sink” simulation from the outset.

For reference, we briefly recapitulate the parameters entering our model (see Gressel et al. 2011). We presume dust with density \( \rho_d = 3 \text{ g cm}^{-3} \) and of a single grain size \( d = 0.1 \mu \text{m} \), which is further assumed to be depleted (by grain growth) to yield a dust-to-gas mass ratio of \( 10^{-3} \). Reactions involving gas-phase ions are represented by magnesium, which is partly bound up in grains and hence taken to be depleted by \( 10^4 \) compared with its solar abundance. In the reaction network, free electrons are created according to the ionization rate \( \xi(\mathbf{x}, t) \), which we evaluate based on the external irradiation. The ionization fluxes are attenuated by the integrated columns of gas to the upper and lower disk surfaces, respectively. The density integration is complicated by the requirement to match the integrals at mesh refinement boundaries, which demands computationally expensive synchronization. To reduce computational overheads, we store \( \eta(\mathbf{r}, t) \) as a passive scalar variable. This allows us to only update \( \eta \) from integrated densities every three computational time steps, which is short enough to reflect potential changes of \( \eta \) on dynamical timescales.

The underlying ionization prescription used here comprises stellar XRs, radioisotopes, and energetic protons and is based on the work of Turner & Drake (2009). We closely follow their approach and restrict ourselves to stellar XRs and interstellar CRs as the prime ionizing agents. The attenuation of interstellar CRs within a PPD has been studied by Umebayashi & Nakano (2009), who derived a dependence

\[
\zeta_{\text{CR}} = 5 \times 10^{-18} \text{s}^{-1} e^{-\Sigma_d/\Sigma_{\text{CR}}} \left[ 1 + \left( \frac{\Sigma_d}{\Sigma_{\text{CR}}} \right) \right]^{-\frac{2}{3}},
\]

with \( \Sigma_{\text{CR}} = 96 \text{ g cm}^{-2} \) being a typical CR attenuation depth (see Umebayashi & Nakano 1981) and with an according contribution from the second gas column \( \Sigma_2 \). As in previous work, we use a simple fit to the Monte Carlo radiative transfer calculations of Igea & Glassgold (1999) and approximate the XR ionization rate via

\[
\zeta_{\text{XR}} = 2.6 \times 10^{-15} \text{s}^{-1} \left[ e^{-\Sigma_d/\Sigma_{\text{XR}}} + e^{-\Sigma_2/\Sigma_{\text{XR}}} \right] \Gamma_{\text{m}}^{-2},
\]

where \( \Sigma_d \) and \( \Sigma_2 \) are the column densities to the upper and lower disk surface, respectively, and where \( \Sigma_{\text{XR}} = 8.0 \text{ g cm}^{-2} \) is a characteristic absorption depth. We finally include an ambient ionization due to the decay of short-lived radioisotopes (SRs). As in previous models (Gressel et al. 2012), the related ionization rate is enhanced \( 10 \times \) compared with the nominal value of \( \zeta_{\text{SR}} = 3.7 \times 10^{-19} \text{s}^{-1} \) (see Turner & Drake 2009). This is done to provide a ceiling on the resistivity so that the associated time-step size does not become prohibitively small. In conclusion, we remark that when a gap is formed in the disk, ideally one should include the radial illumination of the gap edges by the star. For reasons of simplicity, however, we neglect this effect in our current models.

### 3.4. The Accretion Sink

Even with the aid of adaptive mesh refinement, it remains prohibitive to resolve the gaseous envelope accreting onto the planetary core. This is to say that fidelity of resolving small spatial scales is traded off against the ability to cover the (comparatively long) timescales on which gap formation occurs. At the same time, making meaningful predictions about the envelope of hot accreting gas will require proper treatment of radiative effects, which are not currently addressed by our model. Instead focusing on the impact of magnetic fields, we are hence interested in asking the question whether the exterior of the accretion region (i.e., the system comprised of the PPD gap, spiral arms, and the CPD) can, in fact, provide the material into the sphere of influence of the core.

Our accretion sink is similar to that used in Kley (1999) in that a fixed fraction of material is removed from the direct vicinity of the planet each time step. Unlike in earlier work, we specify the timescale for removal as the local free-fall time (of a test particle)

\[
\tau_{\text{ff}} = \frac{\pi}{2} \frac{r_{\text{acc}}^3}{2GM_p},
\]

where \( r_{\text{acc}} \) is defined as 5% of the planet Hill radius—corresponding roughly to Callisto’s semi-major axis. This amounts to a small fraction of the mass within a sphere of \( r_{\text{acc}} \) being removed per time step. To avoid discontinuous behavior in the accretion flow near the planet, removal is weighted with a 3D Gaussian kernel. Owing to limited computational resources, the mass augmented to the planet is artificially enhanced four-fold to facilitate a speed-up of the gap formation process. In cases where we evolve an energy equation, the modification to the gas density is accompanied by a correction in the thermal energy density to keep the temperature constant. Within a sphere of 2\( r_{\text{acc}} \), we apply additional cooling to avoid the build-up of strong pressure gradients that would otherwise modify the accretion flow unphysically and retard the flow of gas toward the sink hole. In this region, the temperature is relaxed toward the initial model on a timescale proportional to the Keplerian angular velocity with respect to the planet potential, resulting in temperatures similar to those observed in the radiation-hydrodynamic simulations of Klahr & Kley (2006). To conclude this section, we note that we do not expect a strong numerical effect from the particular scheme adopted for the accretion sink (see Figure 3 in Machida et al. 2008).

### 4. THE PROTOPLANETARY DISK

Previous studies using both laminar and magnetized-turbulent disks have shown that the presence of an accreting gas giant
planet leads to gap formation and the formation of a CPD surrounding the planet.

The main goals of this paper are to study the effects of magnetic fields and time-dependent ionization levels on the evolution of the gap, the CPD, and the gas-accretion rate onto the planet. To achieve these goals, it is necessary to develop a fiducial HD model for comparison purposes. Coming up with a realistic proxy for a layered turbulent accretion disk is a formidable task in its own right. Clearly, in the presence of a dead zone, a height-dependent form of the disk model described in Section 3. All of the simulations employ three levels of mesh refinement. Each of the rest of the PPD, but eventually this is expected to occur on longer timescales of about a few 100 planet orbits (Nelson et al. 2000), or perhaps longer than this in the presence of a dead zone.

The instantaneous cooling in the isothermal model N2 leads to a higher mass and steeper density profile within the CPD and to a pronounced and tightly wound bi-symmetric spiral, as observed in earlier work (e.g., Lubow et al. 1999). In the adiabatic HD model with prescribed cooling shown in panel (b), the spiral arms are less tightly wound because of warmer temperatures in the inner regions of the CPD. The temperature in the inner regions of the CPD remains at $T \approx 150$ K in the isothermal run N2. Even though rapid cooling is applied in the non-isothermal simulations N2 and M1, temperatures in the inner regions of the CPDs reach $T \approx 2000$ K for these runs (similar to the values $T \approx 1500$ K obtained by Klahr & Kley (2006) in their 3D radiation-hydrodynamic runs). The mass of the CPD is smaller in run N1 and the density profile is shallower (see Figure 10 below). Because each fluid element now has its own thermal history leading to moderate pressure fluctuations in the CPD, the flow around the planet is less regular but still laminar. This changes in the MHD case, which is shown in panel (a), in which the flow pattern becomes turbulent, but otherwise appears similar to the equivalent HD simulation.

In their study of giant planets embedded in fully turbulent disks without dead zones, Nelson & Papaloizou (2003) comment on the fact that the spiral wave structure in turbulent PPDs has a more washed out and diffused appearance than observed in laminar disks. This is clearly not the case here in the midplane because the dead zone contains only relatively low-level density perturbations due to waves propagating from the overlying active layers. The dominant perturbations are the spiral wakes excited by the planet. At higher, more turbulent latitudes in this disk, however, the planet-induced spirals will have a more diffuse appearance. An image of the disk based on
The differences between the isothermal and non-isothermal equation of state become very apparent when looking at poloidal cuts through the simulation domain, presented in Figure 2.

Compared with both the HD run N1, panel (b), and the MHD run, panel (a), the CPD is noticeably thinner for the locally isothermal model, N2, shown in panel (c). All three cases, however, agree in the asymmetry between the inner and outer parts of the CPD, with the inner part of the CPD lying nearest to the central star being vertically thicker than the component lying further away. This feature is very likely due to the thermal model that we adopt, which assumes that gas lying closer to the star is warmer than that further out (this is true for
both the locally isothermal and adiabatic disks with thermal relaxation. In a more realistic model, where the temperature of the CPD is determined by local heating and cooling processes, we suspect that this in–out asymmetry will be less pronounced than displayed by our simulations. We note, however, that in the presence of a gap, a moderately decreasing temperature at the midplane may be expected as one moves out across the gap because reprocessed stellar radiation from the outer gap edge more directly illuminates the inner half of the gap (Turner et al. 2012).

The overall structure of the CPD is quite similar in the non-isothermal HD and MHD cases, indicating that gravity remains the dominant force (the plasma $\beta_{\text{pl}}$, defined as the ratio of thermal to magnetic pressure—typically lies between the values 100 and 1000 in the CPD region; see Figure 13), but it is also clear that the circumplanetary region in the MHD run shows considerably more structure than model N1 due to turbulence. Comparing panels (a) and (b) reveals an inherent limitation in the “enhanced” viscosity prescription that we applied in the HD case: the resolved flow remains laminar and the stochastic nature of the resulting structures observed in the MHD run cannot be captured in this framework. While this may be tolerable when studying the secular evolution of a PPD, it is probably not sufficient for the purpose of studying dynamical processes occurring in CPDs such as satellite formation where the time-dependent structure of the CPD may be crucial.

A low-density funnel along the rotation vector of the CPD can be observed in all three simulations. In the two HD simulations, this funnel is always associated with low angular momentum material falling in onto the poles of the planet. In the MHD run, this is also usually the case, but we have also observed the sporadic occurrence of a (one-sided) magnetically collimated protoplanetary outflow that we will discuss in more detail later in this paper. The sporadic nature appears to be related to the fact that the jet can be disrupted by incoming material accreting through the polar regions.

Taking a global view of the simulation results, we can state that significant differences arise in the non-magnetized simulations when moving from an isothermal equation of state to an adiabatic one with imposed cooling, due to the importance of compressional heating associated with gas flowing into the planet Hill sphere. Magnetic fields also play an important role as they transform the circumplanetary flow from one that is essentially laminar to one that is turbulent and highly fluctuating, described later in Section 6.4.

5. GAP FORMATION

As outlined above, we chose an initial planet mass of 100 $M_{\oplus}$, as this mass puts the planet in the rapid growth phase that we are interested in. It is evident that a combination of tidal torques and gas accretion onto the planet lead to the formation of a significant gap within the PPD. Considering the gap-opening criteria that the tidal torque must overcome the effective viscous torque ($q \gtrsim 40/R_e$, where $q = \dot{M}_p/M_c$ and $R_e$ is the Reynolds number) and the Hill sphere radius of the planet must exceed the local scale height,7 we find that a 100 $M_{\oplus}$ planet just fails to meet the second criterion, but satisfies it when its mass grows above 110 $M_{\oplus}$. The first criterion is fulfilled for planet masses exceeding 60 $M_{\oplus}$ when $\alpha = 2 \times 10^{-3}$ and $h = 0.05$, so the opening and maintenance of a gap is expected in the simulations.

We note that low-mass planets are known to open modest gaps or surface density depressions in disks due to shock dissipation of their spiral wakes (Goodman & Rafikov 2001; Muto et al. 2010), so the above gap-opening criteria should not be interpreted too rigidly.

The gap-opening process is illustrated in Figure 3, where we plot radial profiles of the disk surface density in regular intervals of $\Delta t = 121$ yr, corresponding to roughly 20 orbital times at the planet location. While the gap has not reached a fully stationary state, the deepening of the gap is clearly slowing down toward the end of the simulation. Moreover, there is a trend toward deepening the outer half of the gap, which is also visible in Figure 1. This arises primarily because the disk is cooler just outside of the planet location compared with the corresponding region just interior to it. Waves excited at the outer Lindblad resonances are therefore more nonlinear than their inner Lindblad resonance counterparts. The density contrast between the edge and the gap is roughly 20–30, which is a factor of two to three deeper than reported for the local shearing-box models of Machida et al. (2008). The periodic nature of shearing-box simulations may reduce the depth of the gap, an effect that is avoided in our global simulations. We remark that the simulation referenced above had a similar mass planet ($0.4 M_J \simeq 127 M_{\oplus}$), which was embedded in an inviscid disk without magnetic fields or turbulence.

In Figure 4, we show a comparison of the gap structure between the different models. Because of the different accretion rates (see the discussion in Section 7.1 below), we perform this comparison in two different ways: at a constant planet mass of $M_p = 150 M_{\oplus}$ (upper panel) and at a constant time $t - t_0 = 500$ yr after the insertion of the planet (lower panel).

Gap opening proceeds fastest in the isothermal HD case, N2, presumably due to the more efficient mass growth of the planet in this setup. At a given planet mass, the model M1 shows the deepest gap, followed by the equivalent HD model, N1.

Comparing the simulation results at a fixed evolution time, our non-isothermal MHD and HD runs have quite comparable gap structures (lower panel of Figure 4). There is a noticeable difference in the gap edges, with the MHD model having a higher density inner edge and a lower density outer edge compared with the HD model. While this may have implications for the detailed torque balance experienced by the planet, a detailed examination of this issue is probably not warranted due to the fact that none of the simulations have formed a steady-state gap. Gap opening has been studied extensively in fully MRI-active...
disks. While early studies neglected the vertical stratification (Nelson & Papaloizou 2003; Winters et al. 2003; Papaloizou et al. 2004), this has recently become possible to include (Uribe et al. 2011). These studies generally find that a giant planet embedded in an MRI-turbulent disk opens a wider gap than one embedded in an equivalent HD simulation that utilizes the \( \alpha \) model for viscosity and the upper panel of Figure 4 provides some support for this assertion. We caution, however, that direct comparison with these previously reported trends is difficult because the planet in this paper has a time-dependent mass due to accretion, unlike the previous studies that used a fixed mass.

5.1. Consequences of an Evolving Ionization Fraction

The formation of a gap reduces the column density in run M1, increasing the midplane ionization fraction of disk material there through increased penetration of XRs and CRs. In essence, we observe the dead zone in the vicinity of the planet being ignited into a turbulent state by gap formation. As we have discussed above, gap formation does not run to completion in this simulation, so we find that it is primarily the upper and intermediate layers of the disk in the gap that share common characteristics with the MRI-active surface layers of the rest of the PPD. The gap is deepest in the close vicinity of the planet (i.e., in a wedge of \( \phi \approx \pm \pi/8 \)) and the magnetic Reynolds number, \( R_{m} \equiv c_{s}H/\eta \), has high values in the range \( 10^{2} \text{--} 10^{3} \) there. At the same time, the Ohmic Elsasser number, \( \Lambda \equiv v_{A}^{2}/(\Omega \eta) \), with \( v_{A} \) being the Alfvén speed associated with the vertical field, is well above unity. Accordingly, we find this region in particular to be in a turbulent state all the way down to the midplane. Values of \( \beta_{p} \) are in the typical range for active MRI. As one moves away from the planet position around the orbit, the characteristics of the flow change. Particularly in the midplane and away from the planet, \( R_{m} \) drops to values as low as 100, which is generally insufficient for sustained MRI (Oishi & Mac Low 2011). Moreover, \( \Lambda \) becomes smaller than unity, in agreement with linear theory the flow remains laminar in these regions. At the same time, the plasma there is only very weakly magnetized with \( \beta_{p} \approx 10^{5} \).

Taking a global view of the flow in the PPD and gap region, we can say that gas accretes through the disk toward the planet mainly via the upper active regions. As gas enters the gap, it is pulled down by the star’s gravity and toward the planet by its gravity. The low density there allows strong coupling between the gas and the predominantly azimuthal magnetic field such that the field is advected with the fluid into the gap, where it helps to sustain the MRI. Because our simulations do not allow gap formation to run to completion, gas within the gap region at azimuthal locations away from the planet remains laminar at the midplane. On longer timescales, however, we would expect these regions to also become evacuated such that the MRI could be sustained there also. The end result would be a PPD with active surface layers and a laminar midplane region far from the planet, but a magnetized and turbulent gap in the planet’s vicinity through which material passes as it flows into the planet Hill sphere.

5.2. Consequences of Accreting from a Layered Disk

The feeding of gas into the gap region from the surface layers may have important consequences for the chemistry and dust content of the gas that eventually accretes onto the planet, as dust settling and sequestration in the dead zone may reduce the heavy element content of this gas as it accretes through the disk. The possible implications of this are discussed in Section 8.

An interesting consequence of accretion through a gap in a layered disk is that the gas flow through the surface layers into the gap should be able to continue relatively unimpeded by tidal truncation of the disk even for relatively massive planets, because the magnetic stresses and associated \( \alpha \) values in the upper layers tend to be very large (i.e., \( \alpha > 0.1 \)). As such, these disk layers will not satisfy the viscous gap-opening criterion discussed above and should continue to feed the planet as its mass grows, albeit at a moderate rate because of the low densities in the upper layers. As we discuss later in the paper, the accretion rates observed in the viscous HD runs N1 and N2 show a tendency to continuously decrease as the planet mass increases and the gap deepens. Run M1, however, appears to reach a steady accretion rate that is higher than observed in the laminar runs by the end of the simulation (see Figure 16), apparently for the reason just described.

6. THE CIRCUMPLANETARY DISK

We now discuss the features of the flow in our models on scales appropriate to the planet Hill sphere radius. We begin by describing the horizontal flow features in the midplane, before examining the flow pattern in the meridional plane. We discuss and compare the integrated properties of the CPDs that arise in all of our simulations before focusing on the details of the flow that arise in the MHD simulation M1.

6.1. Horizontal Flow Features in the Midplane

Within the Hill sphere (where the gravitational force of the planet dominates), we observe the formation of a CPD in all simulations as a consequence of the approximate conservation of (specific) angular momentum, \( j \equiv rv_{\phi} \) (with respect to the planet’s position).
In Figure 5, we plot the time-averaged specific angular momentum of the flow region near the planet for the locally isothermal run N2. Superimposed streamlines show that the flow has a high degree of symmetry. Marginally bound trajectories closely follow the Roche lobe (equipotential surfaces through L1 and L2 are shown as black lines), indicating that the flow is only weakly affected by pressure forces. The flow within the CPD is nearly circular at this time, but is observed to become more distorted at late times as discussed in Section 7.2. In comparison, the averaged flow field appears less symmetric in the simulation with an adiabatic equation of state and thermal relaxation toward the initial temperature (the upper panel of Figure 6). Unlike in the previous case, the circulating flow appears to extend beyond the Roche lobe and shows some level of distortion. This is partly due to the fact that the upper panel in Figure 6 corresponds to a later time when gap formation is more developed; we generally find that the CPD flow becomes increasingly distorted as the simulations evolve and gap formation becomes more pronounced. The distortion is apparently also amplified by the enhanced pressure forces present in run N1, combined with the in–out asymmetry that the circumplanetary material displays because of the radial temperature profile imposed on the gas.

This trend of reduced order in the fluid trajectories is enhanced substantially when looking at the MHD case, shown in the lower panel of Figure 6. Here, the flow within the Hill sphere is significantly distorted and less symmetric, presumably due to the combined effect of pressure and magnetic forces. Notably, the flow within the CPD appears to have a non-negligible eccentricity. Magnetic field lines (not shown) are dominantly azimuthal (with respect to the PPD) outside the horseshoe region and moderately compressed in the spiral arms; notably $B_\phi$ has opposite signs in the inner and outer part of the PPD. In the vicinity of the planet orbit (the horseshoe region and just beyond), field lines are generally aligned with the flow, with the exception of the post-spiral shock region where angles up to 90° are obtained. Within the CPD, the field lines follow the winding of the spiral shocks. Overall, the field morphology is consistent with that described by Nelson & Papaloizou (2003).

Figure 5. Time-averaged flow field (arrows) and specific angular momentum (color coded) for a midplane slice of the isothermal HD model. Averages are taken over $\sim 4$ planet orbits after the planet has reached a mass of $150 M_\oplus$. (A color version of this figure is available in the online journal.)

Figure 6. Same as Figure 5, but for the non-isothermal HD model (top panel) and the MHD model (bottom panel). (A color version of this figure is available in the online journal.)

6.2. Poloidal and Three-dimensional Flow Features

During the early phase of evolution, prior to pronounced gap formation, we find that the flow field observed in poloidal slices in the isothermal run N2 is qualitatively very similar to the highly symmetric renderings shown in Figure 5 of Machida et al. (2008) and Figure 4 of Tanigawa et al. (2012). This flow consists of inflow toward the planet at high latitudes and outflow toward the L1 and L2 points near the CPD midplane. During late times, when gap formation is more developed, we find that the flow symmetry breaks down and becomes similar to that displayed by the non-isothermal HD run N1, for which we plot poloidal flow lines near the end of the simulation in Figure 7. Within the outer half of the Hill sphere, material is centrifugally spun-out toward the L2 point near the CPD midplane and then recycled into the vertical accretion flow along the vertical axis of the planet. Unlike at early times and in the (intrinsically symmetric) local simulations of Machida et al. (2008), this is markedly not the case on the side of the CPD facing the star, where we only observe inflow. This difference is also clearly seen in Figures 17 and 18 in Section 7.2 below.

Outside the Hill sphere, the non-isothermal flow shows substantially more complex features than the features seen for an isothermal gas. Most notably, there is a pair of vertically
Figure 7. Projected flow field in the vertical plane connecting the planet with the star. Small vortical return flows within the Hill radius (dotted) are associated with underdense regions already seen in Figure 2. Note the strong vortex pair adjacent to the outer spiral arm.

(A color version of this figure is available in the online journal.)

Figure 8. Backward extrapolation of stream lines passing through the close vicinity of the planet in the final snapshot of the non-isothermal HD model. The size of the cube is 1 au and the left corner points toward the star. For orientation, the gas density is shown. Adjacent to the outer spiral shock, a pair of counter-rotating vortices is visible.

(A color version of this figure is available in the online journal.)

Here). These flow lines are traced from the vicinity of the planet in the CPD and demonstrate that inflow primarily occurs from high latitudes and not along the midplane of the CPD. This feature of the accretion flow is discussed in greater detail in Section 7.2. The flow structure becomes even more tangled up in the MHD case, which is not shown here. A detailed discussion of the intrinsic variability seen in this run is presented in Section 6.4 below.

6.3. Integrated Properties of the CPDs

In Figure 9, we compare time-averaged angular momentum profiles within the CPD midplane of the three main models. Differences due to additional radial pressure support are more pronounced for the part of the disk closer to the planet. This is consistent with the geometrically thicker inner disk seen in the insets of Figure 2. As expected in the absence of strong pressure gradients (that might arise because of compressional heating near the planet), the isothermal model is closest to the Keplerian rotation profile (dashed line). In the MHD model, the CPD rotates the slowest, indicating that magnetic forces, at least at some level, contribute to the overall structure of the disk. Estimating the size of the CPD according to the radius at which the specific angular momentum begins to turn, we find the CPD extends to about half the Hill radius, which is similar to but slightly larger than the CPDs reported in other studies where values of one third of the Hill sphere radius (Ayliffe & Bate 2009a; Machida et al. 2008; Tanigawa et al. 2012) or 0.4rH (Martin & Lubow 2011) have been estimated. These estimates are influenced by the disk temperature because of the role of pressure forces at the CPD outer edge and also by the properties of the inflowing gas such as its angular momentum, so we believe our results are consistent with those obtained in previous studies.

The radial surface density profile of the CPDs from all models is shown in Figure 10. Unlike the two non-isothermal models, the isothermal HD model shows a cusp-like inner disk structure, as well as pronounced spiral features (see Lubow et al. 1999), leading to peaks in the surface density profile. We find the profile is close to a $r^{-0.5}$ dependence, which is shallower than the $r^{-1.5}$ dependence reported by Tanigawa et al. (2012). The reason for this minor discrepancy is likely to be our adoption of a larger sink hole radius. When applying a non-isothermal equation of state, we find lower and nearly constant surface densities and much weaker spiral features due to the higher
temperatures. The inclusion of magnetic field leads to a further reduced disk mass. Generally, the observed surface densities are in the range of the “gas-starved” scenario of Canup & Ward (2002, 2006), which is significantly lower than in the “minimum mass” models of Mosqueira & Estrada (2003). If we integrate the mass contained within the Roche region (black lines in Figures 5 and 6), we obtain $0.041 M_\oplus$ for model M1, $0.063 M_\oplus$ for model N1, and $0.196 M_\oplus$ for model N2. These numbers should be contrasted with the Jovian satellite system mass $\simeq 0.067 M_\oplus$, which, if augmented to solar abundance, would imply the presence of $\simeq 7 M_\oplus$ of gas.

6.4. Time-dependent Morphology in the MHD Case

Having inter-compared the gross properties of the CPD regions for our three runs, we now consider the detailed and time-dependent structure of the CPD region for the magnetized run M1. We postpone a discussion of the geometry of the actual accretion flow onto the planet until Section 7.2. In Figure 11, we present a series of snapshots showing vertical slices of the density, demonstrating the substantial time variability of the accretion flow in this region because of the turbulence in the gap that feeds material into the planet Hill sphere. We see low-density patches within the CPD that are connected with the poloidal vortices that develop within the flow, described in Section 6.2 above. A similar level of temporal variation is seen in the sequence of midplane slices shown in Figure 12.

We plot a vertical slice through the CPD region showing the value of the time-averaged plasma $\beta_p$ value in Figure 13. The CPD is relatively weakly magnetized, with $\beta_p$ typically in the range $10^2-10^3$, while material falling in through the funnel-shaped regions above and below the planet are characterized by equipartition-strength magnetic fields. The field topology in this region can very crudely be described as follows: the large-scale azimuthal field in the MRI-active layers is dragged downward with the accretion flow. Within the Roche lobe, the cusp of the now V-shaped field lines is twisted upward by the circular flow around the planet, producing a helical field topology. One consequence of this field evolution appears to be the launching of a relatively diffuse, accelerated wind from the conical regions above and below the CPD. More localized vertical outflows (e.g., note the mushroom-shaped bow shock seen in the upper half of the fourth panel in Figure 11) appear sporadically during the simulation but generally do not display high levels of collimation and are quickly disrupted by the vigorous MRI turbulence.

In contrast to the rather diffuse disk wind just described, the MHD simulation did produce one episode of sustained, collimated outflow, as shown in the last panel of Figure 11. This is more dramatically illustrated by the volume rendering in Figure 14, which shows both the fluid streamlines and magnetic field topology. This collimated outflow lasted for a period of about four planet orbits before it was squashed by infalling material. Unfortunately, the non-aligned mesh geometry and intrinsic variability make the outflows hard to access quantitatively, so the driving mechanism cannot be pinned down definitively, but it is clear from the twisted field wrapped around the outflowing gas that the protoplanetary jet seen in Figure 14 is collimated by magnetic hoop stresses. Given the importance of magnetic field advection and rotation in the vicinity of the planet, our suggestion that the outflow is magnetocentrifugally driven is well motivated, but speculative.
at the present time. It is possible that field winding generates sufficient magnetic pressure to launch the jet at its base instead, although this is doubtful because field amplification in the CPD is damped efficiently by Ohmic diffusion (see the volume rendering in Figure 14). A deeper analysis of protoplanetary jet launching will be presented in a future publication. We note that CPD jets have been predicted on theoretical grounds by Quillen & Trilling (1998) and Fendt (2003) and they have been observed in local shearing-box simulations by Machida et al. (2006). The physical setup in this latter study was quite different from the one we consider here, however, because we initiate the simulation with a weak magnetic field and strong local fields are built up in the vicinity of the planet by gas accretion and associated advection of magnetic field from the surrounding PPD at late times in the simulation. The study by Machida et al. (2006), however, employed an equipartition-strength vertical magnetic field, leading to the formation of a well-defined bipolar outflow shortly after the simulation was initiated. Although interesting as a phenomenon in its own right, there is no evidence at present that the launching of a CPD jet has a significant influence of the growth of the planet. Examination of the gas accretion rate onto the planet does not indicate that it is influenced significantly during the time of strongly collimated jet launching, implying that momentum and energy injection into the surrounding infalling envelope is not an important process here. A scenario in which wind or jet launching from a CPD may be important for determining the accretion rate onto the planet corresponds to a scaled-down version of the picture presented by Bai & Stone (2013), where the combined effects of Ohmic resistivity and ambipolar diffusion cause angular momentum transport and mass accretion in a PPD to arise through the launching of a magnetized wind. If such a scenario also applies to accretion onto a planet through a CPD, then wind launching will be central to determining the accretion rate onto the planet.

The stochastic nature of the accretion flow into the planet Hill sphere implies that the angular momentum vector of the CPD material should not maintain a constant orientation. This expectation is confirmed by Figure 15, which shows the time evolution of the angle between the angular momentum vector and the vertical axis calculated for all material inside the planet Roche lobe for the various simulations. As expected, we see that both of the HD runs maintain a reasonably constant direction for the CPD angular momentum vector after an initial period of relaxation, with the tilt angle remaining <1°. The MHD run, however, shows high levels of variability over the full time span of the run, with tilt angles reaching up to 15°. Although this measurement applies to all material in the Roche lobe and not
just the material in the well-defined CPD confined to less than half the Hill radius, it seems very likely that the inner part of the CPD close to the planet will experience substantial disturbance that causes it to tilt and precess. In particular, stochastic accretion of material with different specific angular momenta will cause local warping on orbital timescales that may excite bending waves that propagate through the CPD (e.g., Papaloizou & Lin 1995; Larwood et al. 1996), with interesting consequences for waves that propagate through the CPD (e.g., Papaloizou & Lin 1995; Larwood et al. 1996), with interesting consequences for its dynamical evolution and the formation of regular satellite systems in the inner regions. Indeed, we speculate that when all other things are equal, a higher mass planet will host a CPD that is subject to reduced levels of perturbation in the inner satellite forming regions compared with a lower mass planet, by virtue of the strength of the local gravitational field relative to the external perturbing forces that are independent of the planet mass. This may be of relevance when comparing the regular satellite systems of Jupiter and Saturn, where the Jovian system clearly is more substantial than Saturn’s.

At late times during the MHD simulation, the component of the CPD that we model is itself marginally unstable to the MRI—i.e., according to the Elsasser number, \( A \sim 0.2-20 \), based on Ohmic diffusivity alone. We would predict this to be the case from the surface density values displayed earlier in Figure 10. The penetration depth of XRs is \( \sim 9 \) g cm\(^{-2} \) and for CRs it is \( \sim 90 \) g cm\(^{-2} \). Allowing for partial attenuation of these ionizing sources by overlapping material in the gap region, we would predict that the CPD is MRI-active based on the measured column densities (but see Turner et al. 2013, for a more detailed assessment). The magnetic Reynolds number is rather low (\( Rm \sim 100-1000 \)), which implies that a relatively strong vertical net flux will be required to facilitate sustained MRI. Build-up of strong fields through advection into the Hill sphere, and subsequent amplification of the toroidal component by differential rotation, is counteracted by Ohmic diffusion, we might reasonably expect the disk to attain a marginal state of MRI.

Given the stochastic nature of the accretion flow and the associated tilting and warping of the disk, it is hard to assess whether the CPD in our MHD simulations is in fact laminar or MRI turbulent. Within the restrictions of our diffusivity model and based on the associated dimensionless numbers, the latter possibility can certainly not be excluded. If we, on the other hand, estimate the vertical wavelength of the fastest growing (ideal) MRI mode, we find a typical value of \( \lambda_{\text{MRI}} \approx 5 \times 10^{-3} \) au—approximately coinciding with the Nyquist frequency of our finest grid (and hence falling short of the minimal resolution requirement by roughly a factor of five). In view of this issue, and regarding a possible existence of a disk wind and/or magnetocentrifugal jet, even better resolved mesh-refined models are certainly called for.

7. THE ACCRETION FLOW ONTO THE PLANET

As has been pointed out in various previous studies (e.g., Klahr & Kley 2006; Machida et al. 2008; Shabram & Boley 2013), the accretion flow onto the planet in a laminar disk is genuinely 3D. This is also confirmed by all our models, which develop complex flow structures in both the vertical and horizontal directions, with inflow toward the planet arising primarily from high latitudes. In agreement with previous high-resolution studies of the accretion flow onto a giant planet from laminar, inviscid, isothermal disks (Machida et al. 2008; Tanigawa et al. 2012), we also find that material flows away from the planet near the midplane of the CPD and toward the planet from higher latitudes in our locally isothermal run. Outflow away from the planet near the midplane of the CPD is also observed in the two non-isothermal runs, although the situation in these runs is less clear cut.

7.1. Mass Accretion Rates

Before discussing the complex geometry of the accretion flow onto the planet, we consider the total accretion rates that are produced by the simulations. The left panel of Figure 16 shows the evolving planet mass for the three models.\(^8\) The mass accretion rate into the sink particle is plotted in the right panel of the same figure. The inset shows a magnification of the late accretion phase in model M1. The additional (darker) line shows the mass flux through a spherical surface with twice the radius compared with the actual “sink” region, within which gas is removed from the domain. We remind the reader that the constant sink hole radius equals 5% of the Hill sphere radius of the initial planet mass 100 \( M_{\oplus} \). The minimal offset between the curves demonstrates that the mass flux toward the planet is essentially constant near the accretion sink, indicating that the flow of gas into the sink hole is not retarded by any processes occurring in this region. For the magnetized model, after 500 yr,\(^8\) Note that we artificially enhance the mass accumulated by the planet by a factor of four to speed up the tidally induced gap opening. The reported accretion rates, however, reflect the actual gas flow onto the planet and do not include this factor of four enhancement.
we reach a quasi-stationary state, for which we infer a mass accretion rate of \((8.0 \pm 1.4) \times 10^{-3} \, M_\odot \, \text{yr}^{-1}\). In other words, a Saturn-mass planet can grow to Jovian mass within \(\sim 25,000 \, \text{yr}\), a small fraction of the expected disk lifetime. In comparison, the total Maxwell stress measured within the CPD is inferred as \(\alpha \sigma_c \approx 0.01\), translating into an estimated accretion rate of roughly \(10^{-3} \, M_\odot \, \text{yr}^{-1}\), indicating that probably only a fraction of the mass is actually delivered via viscous transport within the disk.

In the other two models, the accretion rate still drops, presumably since gap formation has not completed yet, and there is no MRI-active surface layer with large stresses replenishing material into the gap region. It is evident, however, that contraction of gas onto the planet is more efficient in the locally isothermal model which lacks compressional heating. Even though we apply cooling on the local dynamical timescale to material deep in the planet Hill sphere in the non-isothermal models, compressional heating influences the dynamics significantly through the build-up of a central pressure gradient, with consequent reduction in the accretion rate. As described earlier in Section 4.2, the temperature in the non-isothermal runs reaches values of \(T \approx 2000 \, \text{K}\) in the vicinity of the planet despite the rapid cooling applied there, whereas in the locally isothermal run it remains fixed at \(T \approx 150 \, \text{K}\). It is clear that in addition to consideration of MHD processes, an accurate model of planetary accretion also needs to account for the influence of radiation transport on the infalling material, even during the runaway growth phase of giant planet formation. Nonetheless, the accretion rates between all models vary by less than a factor of two at the end of the simulation when a quasi-steady state has been reached, indicating that quantitatively the results are in agreement regarding rapid gas accretion onto a growing Saturn-mass planet. We conclude that the adoption of two different equations of state induces only a moderate change in the medium-term accretion rates, in agreement with previous analyses (e.g., Ayliffe & Bate 2009b; Machida et al. 2010). In particular, we note that Machida et al. quote accretion rates of \((6–18) \times 10^{-3} \, M_\odot \, \text{yr}^{-1}\) in their discussion, in excellent agreement with our results.

Unlike in the viscous HD runs, the accretion in model M1 is stochastic, showing variation by a factor of two over timescales \(\approx 5 \, \text{yr}\). In terms of the average accretion rate, however, the non-isothermal HD and MHD runs produce very similar values. The MHD run produces a slightly smaller accretion rate during early phases, when magnetic pressure effects within the Hill sphere seem to moderately impede accretion, but at later times when a significant gap has formed, the accretion rate for this run levels off at a higher value than observed for the HD run. As discussed above, this latter effect seems to be due to the large magnetic stresses operating in the MRI-active regions near the PPD surfaces that are not present in a viscous model with constant \(\alpha\). What is clear from these simulations is that magnetic effects within a semi-realistic protostellar disk model do not provide a barrier to planetary gas accretion on scales larger than 5% of the Hill sphere radius. As far as we can tell from our simulations, giant planets can accrete substantial gaseous envelopes within \(\sim 3 \times 10^4 \, \text{yr}\), in broad agreement with earlier studies of non-magnetized disks.

### 7.2. Details of the Accretion Flow Geometry

Given the complex flow structure in the planet Hill sphere reported here and in previous studies, it is interesting to look at the spherical distribution of the mass flux toward the planet and to compare it with the laminar, inviscid, and isothermal models of Tanigawa et al. (2012). This is done in Figure 17, where we look at the mass flux through shells centered around the planet arising in run N1 (see Figure 5 in Tanigawa et al. 2012, note, however, the different set of radii, owing to the lower resolution in our global model compared with their shearing box). Mass fluxes are averaged over approximately one planet orbit and our convention is such that outflow corresponds to positive values (red) and accretion is indicated by negative values (blue). The top four panels display the mass flux onto the planet after \(t \approx 20 \, \text{orbits}\) when its mass equals 150 \(M_\oplus\). These are comparable with Tanigawa et al. (2012) because their model ran for \(\approx 25\) planet orbits, leading only to moderate gap formation. The bottom four panels correspond the end of the simulation when the planet mass moderately exceeds 200 \(M_\oplus\) and gap formation is more pronounced.

Considering the early evolution first, we see that the first panel, at \(r = r_{1\text{H}}\), matches well with the one from Tanigawa et al. (2012). Material flows into the Hill sphere at two diametrically opposite longitudes \(\sim 20^\circ\) in the counter-clockwise direction (when viewed from above) from the L1 and L2 points, respectively. The pronounced \(m = 2\) structure indicates that material also leaves the Hill sphere close to the Lagrange points. The region of inflow seems to be rather loosely defined while the corresponding outflow is more focused, due to the influence of the spiral wake induced by the planet. We remark that the spherical \(r = r_{1\text{H}}\) surface is somewhat misleading in that the true shape of the Roche lobe is closer to a rugby ball than to a football (this is evident in Figure 5). We conclude that the flow pattern seen in the first panel of Figure 17 hence reflects the horseshoe orbits entering and leaving the spherical surface at \(r = r_{1\text{H}}\), which is of little significance when considering accretion onto the planet. For the remaining three panels, we find a qualitatively similar distribution as in Tanigawa et al. (2012). At all radii, the mass fluxes show a clear \(m = 2\) structure related to the spiral features observed in the CPD, again in agreement with Tanigawa et al. (2012). At \(r = 0.5r_{1\text{H}}\), inflow occurs at high latitudes and outflow arises around the midplane region at latitudes between \(\theta' \approx \pm 30^\circ\). On scales of \(r = 0.1r_{1\text{H}}\), there is a noticeable flattening of the disk toward the midplane because the near-constant temperature leads to a flaring CPD structure.

The lower four panels show that the detailed flow geometry toward the planet evolves over time, such that the distinct \(m = 2\) feature observed at all radii during early evolution is lost. At \(r = 0.5r_{1\text{H}}\), we observe a superposition of \(m = 2\) and \(m = 1\) features, and at smaller radii \(m = 1\) dominates. This indicates that the inner CPD develops an elongated/elliptical structure, confirmed by an examination of streamlines in the Hill sphere. The reason for this change in structure is not clear, but we conjecture that a more pronounced openings of the gap at late times may modify the symmetry of the accretion flow onto the CPD arising from the inner and outer disk. This is an effect that can only arise in a global disk model and is not expected in local shearing-box simulations. At late times, the CPD midplane continues to display outflow away from the planet and inflow remains confined to overlying latitudes above and below the midplane. Tanigawa et al. (2012) report that the midplane region continues to display outflow on scales as small as \(r = 0.03\). Taking this result at face value suggests that further simulations are required that probe even deeper into the Hill sphere to examine at which radius the midplane region of the CPD behaves as a classical accretion disk, instead of transporting mass away from the planet.
Figure 17. Angular distribution of mass flux ($\rho v_r$) into spherical shells around the planet (with $r = 1, 0.5, 0.2, \text{and } 0.1$ Hill radii) for the isothermal HD case at early times (top four panels) and at the end of the simulation (bottom four panels). The coordinates ($\theta', \phi'$) = (0, 0) and (0, 180) correspond to the sub-solar and anti-solar points, respectively. Inflow is represented by negative values. The negative contours are represented by dashed lines, whereas positive contours are drawn as solid lines. (A color version of this figure is available in the online journal.)

The situation described above changes only moderately when we consider the two non-isothermal runs, for which the geometry of the mass flux is shown in Figure 18. We note that the isothermal model N2 showed only modest temporal variation over short time intervals, so the plots in Figure 17 show averages taken over one orbital period. As discussed already, both of the non-isothermal models N1 and M1 showed higher levels of temporal variability and Figure 18 displays averages taken over the last 50 orbits of the simulations.

When comparing the last four panels of Figure 17 with Figure 18, we see that the main features of the flow in and out of the Hill sphere at $r = r_H$ are similar in all runs. The basic $m = 1$ symmetry of the mass flux at smaller radii is also similar. The main differences arise because of the flattening of the CPD in model N2, which is not replicated in the hotter CPDs in runs N1 and M1.

Focusing first on the non-magnetized simulation N1, we note that inflow is observed at essentially all longitudes for high latitudes $|\theta'| \gtrsim 30^\circ$. This is true at all radii from $r = 0.5r_H$ down to $r = 0.1r_H$, as seen in Figure 18. We remark that this is consistent with the flow field shown in Figure 7. Considering the flow at lower latitudes in the equatorial region, we see that for $|\theta'| \lesssim 30^\circ$ there lies a contiguous band of inflow at all radii $r \lesssim 0.5r_H$ that lies between $\phi' \approx 330^\circ$ and $\phi' \approx 90^\circ$, straddling the substellar point. Moving around to longitudes lying between $100^\circ \lesssim \phi' \lesssim 330^\circ$, we see a well-defined and outflow near the equator. These bands of inflow and outflow are also consistent with the flow field shown in Figure 7, which shows inflow through the hemisphere facing the star and outflow near the midplane through the hemisphere facing away from the star. Summarizing, we can say that inflow occurs from high latitudes, causing gas to fall toward the planet onto a CPD whose half opening angle is $\approx 30^\circ$. The persistent bands of inflow and outflow at all radii indicate that the flow within the CPD is decidedly non-circular and instead exhibits an elongated or elliptical topology, similar to that observed in run N2 at late times.

Moving on to the magnetized run, we see that the mass flux geometry shown in the last four panels of Figure 18 are very similar to those described for run N1. An examination of individual snapshots during the run shows significant time
dependence on orbital timescales, but the average accretion flow in the magnetized disk looks similar to that in the viscous model.

We now consider the flow rate toward the planet as a function of latitude only. This is done by azimuthally averaging the mass fluxes displayed in Figures 17 and 18. These averaged mass fluxes are displayed in Figure 19. The top panel shows the isothermal model N2 at the end of the simulation and this compares very well with previous results (see Figure 6 in Tanigawa et al. 2012). In particular, we see that mass flow toward the planet at all radii occurs from high latitudes $|\theta| \geq 30^\circ$, with the midplane region showing net outflow away from the planet at these radii. It is interesting to note that the inclusion of viscosity and angular momentum transport in the CPD in our run N2 leads to a net flow profile that agrees qualitatively with the inviscid model presented by Tanigawa et al. (2012). The non-isothermal HD case N1 is shown in the middle panel of Figure 19 and agrees reasonably well with the isothermal run at radii $r = 1$ and $0.5 r_H$, where outflow near the midplane and inflow at high latitudes are observed. At $r = 0.2 r_H$, we see that the mass flux near the midplane is of small magnitude and alternates between inflow and outflow. Net inflow at this radius is dominated by contributions from high latitudes. At $r = 0.1 r_H$, inflow occurs at all latitudes, with a modest contribution from the midplane region and a dominant contribution from higher latitudes. The results from model M1 shown in the bottom panel are broadly consistent with model N1 and again display the important inflow contribution from high latitudes. We again note that the nature of the accretion flow is not dramatically altered by the inclusion of magnetic fields compared with the viscous disk model.

8. SUMMARY AND CONCLUSIONS

In this paper, we have presented the results of global 3D hydrodynamic (HD) and magnetohydrodynamic (MHD) simulations of accreting planets embedded in protoplanetary disks (PPDs). The MHD simulation utilized a detailed ionization model from which Ohmic resistivity was calculated, leading to a disk with active surface layers that sustain MRI turbulence and midplane regions that host a dead zone where MRI turbulence is damped by resistivity. The HD simulations adopted the $\alpha$ model for anomalous viscosity, calibrated to give the same volume-averaged stress as the magnetized disk, and were computed for...
the purpose of comparison with the turbulent disk simulation. One of the HD simulations adopted a locally isothermal equation of state. The MHD run and the other HD simulation both adopted an adiabatic equation of state combined with thermal relaxation, where the temperature in the disk was continuously forced back toward its initial value on the local orbital timescale. All simulations used three levels of adaptive mesh refinement to resolve the region inside the planet Hill sphere. The accretion flow in these models is followed down to a radius equal to 5% of the Hill sphere. Interior to this radius, a sink hole accretes material of the angular momentum vector of material there to the planet. Accretion of gas into the gap and planet Hill sphere occurs from high latitudes. Outflow away from the planet occurs in the midplane regions of the CPD for all radii down to 0.1$r_H$, this being the limit of where we can measure the flow geometry.

2. The locally isothermal HD simulation produces results very similar to those described previously by Machida et al. (2008) and Tanigawa et al. (2012). Infalling material leads to strong inflow toward the planet from high latitudes. Outflow away from the planet occurs in the midplane regions of the CPD for all radii down to 0.1$r_H$, this being the limit of where we can measure the flow geometry.

3. The adoption of an adiabatic equation of state with thermal relaxation in a viscous HD model leads to a warmer and thicker CPD. Even in the presence of rapid cooling, compressional heating increases the temperature to $T \sim 2000 \text{ K}$ in the inner CPD. The increased pressure support reduces the accretion rate onto the planet during early evolution compared with the isothermal model. At late times, when higher angular momentum material accretes into the Hill sphere, the accretion rates converge.

4. In agreement with the locally isothermal run, both the MHD and HD simulations with thermal relaxation demonstrate that accretion toward the planet occurs from high latitudes. The midplane region of the CPD displays net outflow away from the planet down to radii $r \geq 0.2r_H$.

5. Gap opening in the MHD simulation leads to ignition of the dead zone into a turbulent state as XRs and CRs penetrate the gap region. The global structure of the disk is one in which accretion occurs in the active surface layers far from the planet, with the midplane region there being largely inert. Near the planet, there is a deep gap where the ionization fraction allows for the development of MRI turbulence.

6. Enlivening the gap in the MHD simulation causes the accreting planet to be embedded in a turbulent environment. Accretion onto the planet becomes stochastic and the flow in the Hill sphere displays significant temporal variability. The CPD has a surface density $\sim 30 \text{ g cm}^{-2}$ between its outer edge at $\sim 0.5 \text{ Hill radii}$ and the inner boundary at 5% of the Hill radius. This is low enough to be sufficiently ionized by XRs and CRs to sustain the MRI. Our measurement of the Elsasser number there suggests that the CPD should indeed be MRI active, at least in its outer regions.

7. Gas accretion into the gap and planet Hill sphere occurs largely from the active surface layers of the surrounding PPD. As gas enters the gap, it is pulled toward the midplane by the star and planet gravity, dragging the (largely azimuthal) magnetic field with it. Gas and magnetic field lines that enter the Hill sphere join the rotating CPD, leading to the generation of helical magnetic fields. These launch what appear to be sporadic magnetocentrifugally driven outflows. These are generally found to be loosely collimated, but we have observed at least one extended time interval during which a highly collimated jet was launched from the CPD region. The protoplanetary jet does not influence strongly the gas accretion rate onto the planet.

8. Stochastic accretion into the Hill sphere causes the direction of the angular momentum vector of material there to vary significantly. Our model suggests that CPDs display substantial time variability in their tilt angle, amounting to changes $\sim 10^\circ$ on orbital timescales. Although we cannot model the inner regions of the CPD explicitly, we speculate that oscillations in global tilt angle will allow propagation of bending waves into the inner satellite forming region, creating a source of disturbance that may influence the formation of regular satellite systems. In this scenario, it seems likely that satellite building blocks orbiting within

Figure 19. Azimuthally integrated mass fluxes (in code units) as a function of latitude for the isothermal HD model (top), the non-isothermal HD model (middle), and the non-isothermal MHD model (bottom). For low latitudes and $r \geq 0.5r_H$, material is actually expelled (positive values).

(A color version of this figure is available in the online journal.)
the inner CPD will develop mutually inclined orbits if they are strongly coupled to the gaseous CPD, leading to a slower growth of satellites.

9. The accretion rate onto the planet in the MHD simulation reaches a steady value of $M = 8 \times 10^{-3} M_{\oplus} \, \text{yr}^{-1}$, large enough for a Saturn-mass planet to become a Jovian planet in $\sim 3 \times 10^4 \, \text{yr}$, much shorter than the expected disk lifetime. This steady state is reached $\sim 100$ orbits after insertion of the planet. The accretion rates in the viscous HD simulations continue to fall below this value and do not reach steady values by the end of the runs. This is a consequence of the accretion flow in the magnetized disk being confined to the active layers where the magnetic stresses are large. This flow is impervious to tidal truncation and gap formation by the planet, even if its mass becomes quite large, because the effective viscous stress there is large (i.e., $\alpha \gtrsim 0.1$), leading to saturation of the mass accretion rate at a higher value than obtained in a viscous model with a constant $\alpha$ value.

8.1. Implications

The above results have a number of implications for planet and satellite formation. The feeding of gas into the gap region from the surface layers of the PPD may have important consequences for the chemistry and dust content of the gas that eventually accretes onto the planet, as dust settling and sequestration in the dead zone may reduce the heavy element content of this gas as it accretes through the disk. One implication is that gas accreted through a gap in a disk with active surface layers and a midplane dead zone will have lower opacity. Models of giant planet formation (e.g., Pollack et al. 1996; Papaloizou & Nelson 2005; Movshovitz et al. 2010) show that the upper-envelope opacity is crucial in determining the envelope accretion timescale. A low-mass planet (i.e., $5-10 M_{\oplus}$) deeply embedded in a disk without a gap will probably not experience a significant reduction in opacity of the accreted gas as it originates from the local midplane during the early phases of envelope accretion. If placed in a region of the disk where the scale height is much smaller, however, such as in the inner few tenths of an au where $h \lesssim 0.02$ (depending on the viscous dissipation rate, see, e.g., D’Angelo & Marzari 2012), gap formation is expected even for planets with masses $<10 M_{\oplus}$ and these may then accrete metal-poor, low-opacity gas, reducing the core mass required to build up a substantial gaseous envelope. This may provide an explanation for the low-mass and low-density planets that have been discovered by the Kepler mission, such as “Kepler 11e” (Lissauer et al. 2011). One potential caveat is that most MHD models have a transition where the dead zone disappears and the disk becomes fully active at a distance of a few tenths of an au from the central star (Gammie 1996; Iglner & Nelson 2006). We note that a planetary core accreting gas from this disk region will still accrete metal-poor gas, because gas arriving at the inner regions will have accreted through the disk surface layers from further out in the disk where dust settling will have reduced the heavy element content of this gas. An examination of this idea will require the development of detailed planetary envelope accretion models in which the usual outer BC that matches the envelope onto the background protoplanetary nebula model will need to be replaced with one that assumes that the envelope has a quasi-free surface, perhaps surrounded by a thick CPD. Such calculations were presented by Papaloizou & Nelson (2005) for giant planets undergoing runaway gas accretion, but have not been performed for planets in the earlier phase of quasi-static gas settling. Alternatively, this scenario could be explored using 3D radiation-hydrodynamic simulations similar to those presented by Ayliffe & Bate (2009a, 2009b).

The comments above concerning the heavy element content of accreted gas also have implications for the formation of satellite systems. Tanigawa et al. (2012) have noted already that delivery of gas deep into the Hill sphere from high latitudes may lead to accretion of low-metallicity gas onto the inner CPD if grain growth and settling have occurred in the surrounding PPD. The fact that low-metallicity gas is delivered into the gap largely from the active surface layers simply reinforces this point.

The temperatures that we obtained for the inner CPDs in the models with applied cooling should not be taken too seriously given our crude thermal model. In spite of this, it is noteworthy that the inner temperatures obtained of $T \approx 2000 \, \text{K}$ are similar to those reported by Klahr & Kley (2006) in their radiation-hydrodynamic simulations. Temperatures in excess of $T \approx 1500 \, \text{K}$ are large enough to vaporize refractory materials, in addition to volatiles such as water ice, so taken at face value our results and those of Klahr & Kley (2006) indicate that satellite building material is unlikely to be transported to the vicinity of the planet by this early-stage accretion flow. Achieving temperatures low enough to support condensation of ices for building icy satellites clearly requires a phase of evolution in which material is delivered to the inner regions of CPDs at much lower rates because of the dependence of the temperature on the local gas accretion rate.

8.2. Concluding Remarks

Although the model we have presented of a planet accreting gas from a layered PPD represents a significant step forward in terms of complexity, realism, and numerical resolution, there are numerous omissions to the physical model that need to be addressed in future work before we can be confident that gas accretion rates onto forming giant planets are fully understood. The non-ideal MHD model that we have adopted neglects the potentially important effects of ambipolar diffusion and Hall EMFs, both of which may significantly modify the results we have presented (Wardle & Salmeron 2012; Bai & Stone 2013). In a future paper, we will be particularly interested in examining the influence of ambipolar diffusion in the gap region. The low densities there may allow this effect to quench the turbulent nature of the flow in the gap and this will have significant consequences for the time dependence of the flow in the Hill sphere.

Related issues that need to be explored are the transport of stellar XRs into the gap region (as this is crucial for determining the ionization levels close to the giant planet) and the level of CR ionization experienced by the PPD. For XR ionization, we have taken a simple approach and adopted results from the model of Igea & Glassgold (1999) in which XRs are scattered vertically toward the disk midplane from the overlying disk atmosphere, with attenuation depending simply on the column density. This model, however, does not account for the effects of gap opening and a more accurate approach will require transfer calculations of XRs being scattered into the gap. A further refinement would be to include the effects of XR flares. These lead to both substantial increases in the XR flux and also to a substantial hardening of the spectra (Arzner et al. 2007; Fracisiosini et al. 2007; Getman et al. 2008). Given that hard XRs are able to penetrate more deeply into the disk, this may influence the ionization of material in the vicinity of the planet. CR ionization rates experienced at PPD surfaces and the attenuating influences
of magnetized stellar winds and disk winds were examined recently by Cleeves et al. (2013). By extrapolating from the solar wind modulation of CRs in the heliosphere, with its observed dependence on sunspot coverage, these authors conclude that the stellar wind can reduce CR ionization rates at PPD surfaces by many orders of magnitude compared with the value adopted in this paper. The disk wind and its associated ordered field, in contrast, have little influence because of the compensating effects of magnetic funnelling and mirroring. An issue that remains to be addressed is the interaction between the stellar wind and the disk wind. If it carries sufficient momentum, the disk wind may retard the equatorial flow of the stellar wind and confine it to polar regions, allowing significant penetration of CRs to the disk surfaces. If the disk wind is ineffective in this regard, however, then it seems likely that CR ionization rates will indeed be low. The effect of this on our results will be to simply reduce the depths of the active layers in the PPD but not remove them altogether because the stellar XR\`es penetrate to column densities of \( \sim 10^6 \text{ g cm}^{-2} \). The enlivening of the gap region will still arise because the column densities there become low enough for significant penetration by XR\`es.

A proper treatment of the gas thermodynamics is also needed. This is illustrated by the fact that during the time before a steady state has been achieved, our locally isothermal model displays a higher gas accretion rate than do the models with an adiabatic equation of state and local cooling. In other words, thermodynamics play a role in regulating the accretion rate. When a steady state has been achieved, our locally isothermal model displays a higher gas accretion rate than do the models with an adiabatic equation of state and local cooling. In other words, thermodynamics play a role in regulating the accretion rate. When a steady state has been achieved, our locally isothermal model displays a higher gas accretion rate than do the models with an adiabatic equation of state and local cooling. In other words, thermodynamics play a role in regulating the accretion rate.

We thank the anonymous referee for useful comments that led to an improvement of this paper. Part of the research was carried out at the Jet Propulsion Laboratory, California Institute of Technology, under a contract with the National Aeronautics and Space Administration. The simulations presented in this paper were run on the QMUL HPC facility. Three-dimensional imagery produced by vapor (Clyne et al. 2007; www.vapor.ucar.edu), a product of the Computational Information Systems Laboratory at the National Center for Atmospheric Research.

**APPENDIX**

It is clear that insufficient numerical resolution will lead to erroneous results being obtained in terms of stresses and accretion rates. For example, Fromang & Nelson (2006) presented one of the first systematic studies of MRI-driven turbulence in vertically stratified global disk models and found that approximately six cells were required to resolve the fastest growing MRI modes in order for nonlinear MRI turbulence to be sustained in disks containing weak net-toroidal magnetic fields. In recent work, the influence of numerical resolution in producing robust results has been emphasized strongly (Hawley et al. 2013), where it has been suggested that between 20 and 30 cells per unstable MRI wavelength are required to obtain reliable results in disk models that sustain MRI turbulence throughout (i.e., no dead zone is present).

The basic disk model, M1, presented in the main body of this paper, has a significantly lower resolution than this and so we have computed a suite of test calculations to examine the influence of changing the numerical resolution in the simulations. These calculations did not use mesh refinement. The size of the azimuthal domain used was \( \pi / 4 \) instead of \( \pi / 2 \), but we adopted the same grid size ratio \( \Delta r : \Delta \phi \) as in the original simulation. The time evolution of the Maxwell stress (as defined in Equation (13) of Fromang & Nelson 2006) evaluated between the radii 3 \( \leq r \leq 4 \) au for the different models is shown in the left-hand panel of Figure 20. The number shown in the inset is the number of radial grid cells, \( N_r \), where \( N_r = 384 \) in model M1. When changing the radial resolution, we also changed the resolution in the other dimensions by the same factor. We see that there is little evidence of a systematic trend in the results as a function of numerical resolution. This point is further supported by the right-hand panel of the same figure, where we show radial profiles of the temporally averaged Maxwell stresses. Apart from the low-resolution model with 192 grid cells in radius, all models produce nearly indistinguishable results.

The reasons for this are two-fold. First, we adopt a net-vertical magnetic field in our simulations that is essentially conserved throughout the runs. This has the effect of reducing the influence of the numerical resolution on the results because we do not rely on any sort of disk dynamo to maintain the large-scale vertical field. Second, the presence of the dead zone means that only a relatively narrow region of the disk above the midplane is susceptible to growth of the MRI—there the fields are generally quite strong (compared with the gas pressure), leading to MRI developing on comparatively larger scales. In contrast, the resolution requirements are more severe in a disk where turbulence can also develop near the midplane because the value of \( \beta_p \equiv P_{\text{gas}} / P_{\text{mag}} \) will be larger there. The low density and pressure in the upper disk layers means that \( \beta_p \) is fairly small (between 20–30) and this allows the MRI to be quite well resolved even for weak fields. For ideal MHD, the wavelength of the fastest growing mode can be approximated by

\[
\lambda_{\text{MRI}} \approx 2 \pi \sqrt{v_s / \Omega},
\]

where \( c_s \) and \( H \) are the sound speed and vertical scale height. In the context of this paper, the sound speed is expressed in terms of the gas pressure, which is \( P_{\text{gas}} \). We can express this as

\[
\lambda_{\text{MRI}} \approx 2 \pi H / \sqrt{\beta_p}.
\]

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It is clear that insufficient numerical resolution will lead to erroneous results being obtained in terms of stresses and accretion rates. However, Fromang & Nelson (2006) presented one of the first systematic studies of MRI-driven turbulence in vertically stratified global disk models and found that approximately six cells were required to resolve the fastest growing MRI modes in order for nonlinear MRI turbulence to be sustained in disks containing weak net-toroidal magnetic fields. In recent work, the influence of numerical resolution in producing robust results has been emphasized strongly (Hawley et al. 2013), where it has been suggested that between 20 and 30 cells per unstable MRI wavelength are required to obtain reliable results in disk models that sustain MRI turbulence throughout (i.e., no dead zone is present).

The basic disk model, M1, presented in the main body of this paper, has a significantly lower resolution than this and so we have computed a suite of test calculations to examine the influence of changing the numerical resolution in the simulations. These calculations did not use mesh refinement. The size of the azimuthal domain used was \( \pi / 4 \) instead of \( \pi / 2 \), but we adopted the same grid size ratio \( \Delta r : \Delta \phi \) as in the original simulation. The time evolution of the Maxwell stress (as defined in Equation (13) of Fromang & Nelson 2006) evaluated between the radii 3 \( \leq r \leq 4 \) au for the different models is shown in the left-hand panel of Figure 20. The number shown in the inset is the number of radial grid cells, \( N_r \), where \( N_r = 384 \) in model M1. When changing the radial resolution, we also changed the resolution in the other dimensions by the same factor. We see that there is little evidence of a systematic trend in the results as a function of numerical resolution. This point is further supported by the right-hand panel of the same figure, where we show radial profiles of the temporally averaged Maxwell stresses. Apart from the low-resolution model with 192 grid cells in radius, all models produce nearly indistinguishable results.

The reasons for this are two-fold. First, we adopt a net-vertical magnetic field in our simulations that is essentially conserved throughout the runs. This has the effect of reducing the influence of the numerical resolution on the results because we do not rely on any sort of disk dynamo to maintain the large-scale vertical field. Second, the presence of the dead zone means that only a relatively narrow region of the disk above the midplane is susceptible to growth of the MRI—there the fields are generally quite strong (compared with the gas pressure), leading to MRI developing on comparatively larger scales. In contrast, the resolution requirements are more severe in a disk where turbulence can also develop near the midplane because the value of \( \beta_p \equiv P_{\text{gas}} / P_{\text{mag}} \) will be larger there. The low density and pressure in the upper disk layers means that \( \beta_p \) is fairly small (between 20–30) and this allows the MRI to be quite well resolved even for weak fields. For ideal MHD, the wavelength of the fastest growing mode can be approximated by

\[
\lambda_{\text{MRI}} \approx 2 \pi \sqrt{v_s / \Omega},
\]

where \( c_s \) and \( H \) are the sound speed and vertical scale height. In the context of this paper, the sound speed is expressed in terms of the gas pressure, which is \( P_{\text{gas}} \). We can express this as

\[
\lambda_{\text{MRI}} \approx 2 \pi H / \sqrt{\beta_p}.
\]
this relation, we estimate that in our case $\lambda_{\text{MRI}}$ is on the order of the pressure scale height itself and thus adequately resolved.

Another issue of potential importance is the role of mesh refinement in determining the evolution of the MRI in the disk and the strength of the turbulent stresses. In a disk without a dead zone, where MRI turbulence can develop throughout the body of the disk, we would expect the region in which mesh refinement is applied to show enhanced activity compared with the surrounding disk regions. This is because shorter-wavelength MRI modes will be resolved in this region, potentially boosting the strength of the turbulence there. In our model, however, the mesh refinement is applied in a spherical region centered around the planet near the midplane. The refined region therefore lies almost entirely within the dead zone and has essentially no effect on the turbulence in the active zones near the disk surfaces. We have tested this by running a simulation with refinement applied around the nominal planet location, but with the planet’s gravity switched off. When the planet gravity is switched on, however, gap formation ensues and we might be concerned that the developing turbulence in the annular region centered on the planet semi-major axis will be affected by the refinement. We note, however, that the refinement is applied in a spherical region around the planet and therefore does not extend far in azimuth. As such, we conclude that the refinement has little influence on the turbulence in the surrounding PPMD and is only effective in the region where the planet gravity dominates, allowing small-scale fluid and magnetic structures to be resolved there, as intended.

REFERENCES

Alexiades, V., Amiez, G., & Gremaud, P.-A. 1996, CNME, 12, 31
Arzner, K., Güdel, M., Briggs, K., Telleschi, A., & Audard, M. 2007, A&A, 468, 477
Ayliffe, B. A., & Bate, M. R. 2009a, MNRAS, 397, 657
Ayliffe, B. A., & Bate, M. R. 2009b, MNRAS, 393, 49
Bai, X.-N., & Stone, J. M. 2013, ApJ, 769, 76
Balbus, S. A., & Hawley, J. F. 1998, RvMP, 70, 1
Balsara, D. S., & Meyer, C. 2010, arXiv:1003.0018
Balsara, D. S., & Spicer, D. S. 1999, JCoPh, 149, 270
Boss, A. P. 1998, ApJ, 503, 923
Bryden, G., Chen, X., Lin, D. N. C., Nelson, R. P., & Papaloizou, J. C. B. 1999, ApJ, 514, 344
Canup, R. M., & Ward, W. R. 2002, AJ, 124, 3404
Canup, R. M., & Ward, W. R. 2006, Natur, 441, 834
Cleeves, L. I., Adams, F. C., & Bergin, E. A. 2013, ApJ, 772, 5
Clyne, J., Mininni, P., Norton, A., & Rast, M. 2007, NJPh, 9, 301
D'Angelo, G., Henning, T., & Kley, W. 2003, ApJ, 599, 548
D'Angelo, G., & Marzari, F. 2012, ApJ, 757, 50
De Val-Borro, M., Edgar, R. G., Artymowicz, P., et al. 2006, MNRAS, 370, 529
Fendt, C. 2003, A&A, 411, 623
Flock, M., Dzyurkevich, N., Klahr, H., & Mignone, A. 2010, A&A, 516, A26
Franciosini, E., Pili, L., Stelzer, B., et al. 2007, A&A, 468, 485
Fromang, S., & Nelson, R. P. 2006, A&A, 457, 343
Gammie, C. F. 1996, ApJ, 457, 355
Gardinier, T. A., & Stone, J. M. 2008, JCoPh, 227, 4123
Getman, K. V., Feigelson, E. D., Micela, G., et al. 2008, ApJ, 688, 437
Goodman, J., & Rafikov, R. R. 2001, ApJ, 552, 793
Gressel, O., Nelson, R. P., & Turner, N. J. 2011, MNRAS, 415, 3291
Gressel, O., Nelson, R. P., & Turner, N. J. 2012, MNRAS, 422, 1140
Hawley, J. F., Richers, S. A., Guan, X., & Krolik, J. H. 2013, ApJ, 772, 102
Hayashi, C. 1981, PThPS, 70, 35
Igea, J., & Glassgold, A. E. 1999, ApJ, 518, 848
Ilgner, M., & Nelson, R. P. 2006, A&A, 445, 205
Klahr, H., & Kley, W. 2004, A&A, 445, 747
Kley, W. 1999, MNRAS, 303, 696
Larwood, J. D., Nelson, R. P., Papaloizou, J. C. B., & Terquem, C. 1996, MNRAS, 282, 597
Lin, D. N. C., & Papaloizou, J. 1986, ApJ, 309, 846
Lin, D. N. C., & Papaloizou, J. C. B. 1993, in Protostars and Planets III, ed. E. H. Levy & J. I. Lunine (Tucson, AZ: Univ. Arizona Press), 749
Lissauer, J. J., Fabrycky, D. C., Ford, E. B., et al. 2011, Natur, 470, 53
Londrillo, P., & del Zanna, L. 2004, JCP, 195, 17
Lubow, S. H., Seibert, M., & Artymowicz, P. 1999, ApJ, 526, 1001
Machida, M. N., Inutsuka, S.-i., & Matsumoto, T. 2006, ApJL, 645, L129
Machida, M. N., Kokubo, E., Inutsuka, S.-i., & Matsumoto, T. 2008, ApJ, 685, 1220
Machida, M. N., Kokubo, E., Inutsuka, S.-i., & Matsumoto, T. 2010, MNRAS, 405, 1227
Martin, R. G., & lubow, S. H. 2011, MNRAS, 413, 1447
Martin, R. G., & Lubow, S. H. 2013, MNRAS, 432, 1616
Meyer, C. D., Balsara, D. S., & Aslam, T. D. 2012, MNRAS, 422, 2102
Miyoshi, T., & Kusano, K. 2005, JCoPh, 208, 315
Muto, T., Suzuki, T. K., & Inutsuka, S.-i. 2010, ApJ, 724, 448
Movshovitz, N., Bodenheimer, P., Podolak, M., & Lissauer, J. J. 2010, Icar, 209, 616
Muto, T., Suzuki, T. K., & Inutsuka, S.-i. 2010, ApJ, 724, 448
Nelson, R. P., Gressel, O., & Umurhan, O. M. 2013, MNRAS, 435, 2610
Nelson, R. P., & Papaloizou, J. C. B. 2003, MNRAS, 339, 993
Nelson, R. P., Papaloizou, J. C. B., Masset, F., & Kley, W. 2000, MNRAS, 318, 18
Oishi, J. S., & Mac Low, M.-M. 2011, ApJ, 740, 18
Okaizumi, S., & Hirose, S. 2011, ApJ, 742, 65
Paardekooper, S.-J., & Mellema, G. 2008, A&A, 478, 245
Papaloizou, J. C. B., & Lin, D. N. C. 1995, ARA&A, 33, 505

Figure 20. Resolution study of the unperturbed disk model in the absence of mesh refinement. Numbers are for the adopted number of radial grid cells; the resolution in the vertical and azimuthal directions are scaled accordingly. Left: time series of the volume-averaged Maxwell stress evaluated in the radial range $3 \leq r \leq 4$ au. Right: radial profile of the volume- and time-averaged Maxwell stress. (A color version of this figure is available in the online journal.)
Papaloizou, J. C. B., & Nelson, R. P. 2005, A&A, 433, 247
Papaloizou, J. C. B., Nelson, R. P., & Snellgrove, M. D. 2004, MNRAS, 350, 829
Pierens, A., & Nelson, R. P. 2010, A&A, 520, A14
Pinte, C., Padgett, D. L., Ménard, F., et al. 2008, A&A, 489, 633
Pollack, J. B., Hubickyj, O., Bodenheimer, P., et al. 1996, Icar, 124, 62
Quillen, A. C., & Trilling, D. E. 1998, ApJ, 508, 707
Sano, T., Miyama, S. M., Umebayashi, T., & Nakano, T. 2000, ApJ, 543, 486
Saumon, D., & Guillot, T. 2004, ApJ, 609, 1170
Shabram, M., & Boley, A. C. 2013, ApJ, 767, 63
Skinner, M. A., & Ostriker, E. C. 2010, ApJS, 188, 290
Tanigawa, T., Ohtsuki, K., & Machida, M. N. 2012, ApJ, 747, 47
Turner, N. J., Choukroun, M., Castillo-Rogez, J., & Bryden, G. 2012, ApJ, 748, 92

Turner, N. J., & Drake, J. F. 2009, ApJ, 703, 2152
Turner, N. J., Lee, M. H., & Sano, T. 2013, arXiv:1306.2276
Umebayashi, T., & Nakano, T. 1981, PASJ, 33, 617
Umebayashi, T., & Nakano, T. 2009, ApJ, 690, 69
Uribe, A. L., Klahr, H., Flock, M., & Henning, T. 2011, ApJ, 736, 85
Ward, W. R. 1997, Icar, 126, 261
Wardle, M. 2007, Ap&SS, 311, 35
Wardle, M., & Salmeron, R. 2012, MNRAS, 422, 2737
Watson, A. M., Stapelfeldt, K. R., Wood, K., & Ménard, F. 2007, in Protostars and Planets V, ed. B. Reipurth, D. Jewitt, & K. Keil (Tucson: AZ: Univ. Arizona Press), 523
Winters, W. E., Balbus, S. A., & Hawley, J. F. 2003, ApJ, 589, 543
Ziegler, U. 2004, JCoPh, 196, 393
Ziegler, U. 2011, JCoPh, 230, 1035
Zingale, M., Dursi, L. J., ZuHone, J., et al. 2002, ApJS, 143, 539