Orbital Evolution of Moons in Weakly Accreting Circumplanetary Disks

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Abstract

We investigate the formation of hot and massive circumplanetary disks (CPDs) and the orbital evolution of satellites formed in these disks. Because of the comparatively small size-scale of the subdisk, quick magnetic diffusion prevents the magnetorotational instability (MRI) from being well developed at ionization levels that would allow MRI in the parent protoplanetary disk. In the absence of significant angular momentum transport, continuous mass supply from the parental protoplanetary disk leads to the formation of a massive CPD. We have developed an evolutionary model for this scenario and have estimated the orbital evolution of satellites within the disk. We find, in a certain temperature range, that inward migration of a satellite can be stopped by a change in the structure due to the opacity transitions. Moreover, by capturing second and third migrating satellites in mean motion resonances, a compact system in Laplace resonance can be formed in our disk models.

Key words: planets and satellites: formation – planets and satellites: gaseous planets – protoplanetary disks

1. Introduction

Regular satellites around gas-giant planets are thought to form in a surrounding gaseous disk. This notion is supported by the near-circular orbits of the moon systems in our own solar system, which are well-aligned to their equatorial planes, except for irregular satellites that have been captured dynamically.

In analogy with a minimum mass solar nebula (Hayashi 1981) for planet formation, minimum mass sub-nebula models were introduced for satellite formation (e.g., Lunine & Stevenson 1982). For explaining certain characteristics of the Galilean moons, Canup & Ward (2002, 2006) introduced the so-called gas-starved disk model and reproduced the ratio between the planet mass and the total mass of the satellite system. Satellite formation around gas giants in a solids-enhanced minimum mass model was discussed by Mosqueira & Estrada (2003a, 2003b); moreover, Estrada & Mosqueira (2006) developed a scenario in gas-poor environment. Sasaki et al. (2010) introduced an inner cavity in a gas-starved disk and found that several moons in a resonance can be formed. Based on this work, Oghara & Ida (2012) performed N-body simulations and have shown that moons are commonly captured in a 2:1 mean motion resonance outside the cavity and Galilean-like configuration can be formed.

In contrast to the nebula hypothesis, relatively small rocky satellites present today can be well explained by a formation scenario in a tidal spreading disk (Charnoz et al. 2010; Crida & Charnoz 2012; Hyodo et al. 2015). The model is compelling, but it requires a disk of solid material as the starting point. Such circumplanetary “debris" disks may originate from the capture of planetesimals (e.g., Hyodo et al. 2016, 2017) or from tidal disruption of a previous generation of satellites. However, larger satellites, especially those that maintain an atmosphere, need gas around them during their accretion. Therefore, it is reasonable to assume that at least some of the regular satellite systems must have originated from gaseous circumplanetary disks (CPDs).

When a protoplanet grows to the size of several Earth masses in a protoplanetary disk (PPD) comparable to the one typically assumed for the early solar nebula (Hayashi 1981), gas around the planet starts to accrete onto it. At that time, because of the conservation of angular momentum, a rotationally supported disk forms around the planet. On theoretical grounds, CPDs can be observed in many hydrodynamic and magnetohydrodynamic (MHD) simulations (e.g., Tanigawa & Watanabe 2002; Klahr & Kley 2006; Ayliffe & Bate 2009; Machida et al. 2010; Gressel et al. 2013; Szulágyi et al. 2014, 2017; Perez et al. 2015). Because of the orders of magnitude difference in spatial scales, however, resolving the very vicinity of the planet in those simulations is still difficult. Tanigawa et al. (2012) have successfully measured the mass infall rate onto a CPD during the early stage of its evolution, but the long-term evolution remains to be established. Yet, modeling a CPD fully self-consistently during the full PPD and planetary gap evolution is difficult just like modeling the formation of PPDs from cloud-collapse is difficult.

In PPDs, the magnetorotational instability (MRI) is thought to play an important role in facilitating the accretion of the disk gas. Although angular momentum transfer in a CPD was previously expected to be as effective as that in a PPD, sustaining the MRI is more difficult in CPDs (Fujii et al. 2011, 2014; Keith & Wardle 2014; Turner et al. 2014). Thermal ionization can trigger MRI at the inner radii of a CPD if the temperature becomes sufficiently high (Keith & Wardle 2014). However, in the absence of strong MRI turbulence, gas accretion may not be efficient enough to prevent a CPD from becoming massive by accumulation of infalling material. An alternative scenario for angular momentum transport within the CPD may be provided by a magnetocentrifugal disk wind that has been found to operate sporadically in resistive-MHD simulations (Gressel et al. 2013). It remains to be shown whether CPD winds are equally emerging when including additional micro-physics such as ambipolar diffusion. In any case, disk outflows are generally competing with infall, and it is unclear how a steady state can be reached for CPDs that are still deeply embedded in their parent disks.

If the subdisk grows so massive as to become gravitationally unstable, spiral arms appear and transfer angular
momentum. Whether or not the gravitational energy is converted into heat in situ is still under debate, but supposing the energy deposition is local, the temperature with the CPD can become high. In such a situation, episodic accretion caused by a combination of the gravitational instability (GI) and the MRI (boosted by thermal ionization) is to be expected. Martin & Lubow (2011, 2013) and Lubow & Martin (2012, 2013) studied this phenomenon in a layered CPD model that is developed in the context of PPDs (e.g., Armitage et al. 2001).

As mentioned earlier, a CPD is likely to become massive in the absence of significant transport of angular momentum, that is, if the temperature is insufficient to maintain MRI turbulence. In this paper, we develop an alternative model of the satellite-forming region of CPDs considering the mass inflow from the PPD as the dominant factor. Because there still is a gap inside a few tens of planet radii forming region of CPDs considering the mass in

In this paper, we develop an alternative model of the satellite-forming region of CPDs—through numerical simulations.

By doing so, in some of our models, we find a bump in the radial surface-density structure. We will examine whether such a specific location can stop the migration of moons. Given the situation that the innermost moon survives rapid inward type-I migration by convergent migration to the pressure maximum, we investigate the possibility of trapping the second and third moons in a mean motion resonance (MMR). The inner three of the Galilean moons are known to be in a 4:2:1 mean motion resonance, a so-called Laplace resonance. In some of our models, we successfully obtained a system in Laplace resonance.

This paper is organized as follows. In Section 2, we describe our sub-disk model and assumptions, and resulting disks are shown in 3. We then highlight several models with interesting structure and discuss the orbital evolution of moons in the disks in Section 4. Discussion of the obtained results and a brief summary are given in Sections 5 and 6, respectively.

### 2. Modeling of Circumplanetary Disks

#### 2.1. Derivation of Surface Density and Temperature

The equation for the time evolution of surface density is essentially derived in the same way as in Fujii et al. (2014), but in addition, here we simultaneously solve for the temperature structure of the embedded sub-disk.

We determine the surface-density profile of the CPD by solving a diffusion equation with an additional source term stemming from mass infall from the parent PPD. When the sub-disk’s angular velocity is taken to be Keplerian, the evolution of the surface density is given by

\[
\frac{\partial \Sigma}{\partial t} = \frac{1}{r} \frac{\partial}{\partial r} \left[ 3r^{1/2} \frac{\partial}{\partial r} \left( r^{1/2} \nu \Sigma \right) \right] + \frac{f}{r},
\]

where \( r \) is the radius, \( \nu \) is the kinematic viscosity coefficient, and \( f \) is the mass flux from infall onto the CPD. We employ the standard \( \alpha \) prescription of Shakura & Sunyaev (1973), namely,

\[
\nu = \alpha c_s^2 / \Omega_K,
\]

with \( c_s \) being the sound speed and \( \Omega_K \) the Keplerian rotation frequency.

To determine a prescription for the source term \( f \), we adopted the results of a detailed analysis of the 3D high-resolution simulations by Tanigawa et al. (2012). Even though they employed 11 levels of nested grids, the resolution is insufficient to resolve the vertical structure of the CPD in the innermost several Jupiter radii from the planet. Therefore, Tanigawa et al. measured physical values of infalling material at high altitude, where the infall is supersonic—the idea being that the infall rate obtained in this way is not affected by the uncertainty caused by the architecture of the CPD further downstream. The effective mass flux onto an inner part of CPD is

\[
f(r) \propto r^{-1} \quad \text{(see Figure 15 of Tanigawa et al. 2012)}.
\]

We assume that the planet has 0.4 \( M_J \) and lies 5.2 au away from a solar mass star. Based on the values at this distance in the minimum mass solar nebula (Hayashi 1981), the local surface density and sound velocity of the PPD are

\[
\begin{align*}
\varepsilon & = 143 \text{ g cm}^{-2} \text{ and } 4.58 \times 10^3 \text{ cm s}^{-1}, \quad \text{respectively, where } \varepsilon \text{ is a scaling factor representing the reduction of the surface density due to gas dissipation. With these values, we obtain the mass infall rate as}
\end{align*}
\]

\[
f = \begin{cases} 
1.3 \times 10^{-3} \varepsilon \left( \frac{r}{R_J} \right)^{-1} \text{ g cm}^{-2} \text{ s}^{-1} & \text{if } r \leq 20 R_J, \\
0 & \text{if } r > 20 R_J,
\end{cases}
\]

where \( R_J \) is 1 Jupiter radius. As the power-law index of the mass infall rate drops outside ~0.4 Hill radii from the planet (Tanigawa et al. 2012), we set \( f = 0 \) at \( r > 20 R_J \). The initial value \( \varepsilon = 1 \) corresponds to the beginning of the mass infall, and a smaller value indicates a smaller mass infall rate. Since the viscous timescale of the CPD is sufficiently smaller than that of the PPD, we treat \( \varepsilon \) as a constant here.

We adopt the following simplified form used in previous studies (Cannizzo 1993; Armitage et al. 2001) to estimate the temperature structure, that is,

\[
\frac{\partial T_c}{\partial t} = \frac{2(Q_+ - Q_-)}{c_p \Sigma} - v_r \frac{\partial T_c}{\partial r},
\]

where \( Q_+ = (9/8) \nu \Sigma \Omega_K^2 \) represents the viscous heating, \( Q_- = \sigma (1 + 3/8 \kappa \Sigma) T^4_d \) is the radiative cooling, \( \sigma \) is the Stefan–Boltzmann constant, and \( v_r \) is the radial velocity. Further to this, the opacity \( \kappa \) is given by Bell & Lin (1994), summarized in Table 1 (see also Table 1 of Kimura & Tsuribe 2012) as

\[
\kappa = \kappa_0 \rho^\alpha T^b,
\]

where \( \rho \) is the density and \( T \) is the temperature. The specific heat is given by \( c_p \approx 2.7 \mathcal{R} / \mu \), where \( \mathcal{R} \) is the ideal gas constant and \( \mu = 2.34 \) is the mean molecular weight.
We solve Equations (1) and (3) numerically with boundary conditions of zero torque and vanishing temperature gradient at the inner and outer boundaries. The temperature of the PPD at the location of the sub-disk is assumed to be \( T = 123 \) K. We set the temperature of the CPD to this value whenever the calculated temperature is lower than this.

### 2.2. Origin of Viscosity

In this section, we explain how to obtain an estimate for the kinematic viscosity coefficient. The best-known origin of effective viscosity in accretion disks is via MRI turbulence. There are two criteria for the MRI that must be fulfilled for the instability to be active: the disk gas must be ionized enough to be coupled with magnetic field, and the magnetic field is not too strong (Balbus & Hawley 1991; Sano & Miyama 1999; Okuzumi & Ormel 2013). We consider the MRI if the following two conditions are satisfied:

1. the Elsässer number, \( \Lambda = \frac{v_a^2}{\eta \Omega_K} > 1 \), with \( v_a = |B_\phi|/\sqrt{4\pi \rho} \) being the vertical component of the Alfvén velocity, and where \( \eta \) is the magnetic diffusivity, and
2. the \( z \) component of plasma beta defined in terms of net magnetic flux, \( \beta_z = 2c_s^2/v_a^2 \gg 2000 \); see Okuzumi & Ormel (2013) and Fujii et al. (2014) for details.

As suggested by the results of Fujii et al. (2011, 2014), Turner et al. (2014), and Keith & Wardle (2014), sub-disks are not likely to widely sustain well-developed MRI turbulence in the absence of thermal ionization.

If the disk is not subject to sustaining MRI turbulence, material may pile up and the disk eventually becomes gravitationally unstable. Accordingly, if the Toomre parameter,

\[ Q \equiv \frac{\Omega_K c_s}{\pi G \Sigma}, \]

(with \( G \) the gravitational constant) becomes smaller than a value of 2, we employ an effective viscosity of \( \alpha_{GI} \equiv \exp(Q^{-4}) \) (Zhu et al. 2010; Takahashi et al. 2013). In such a case, the disk can be easily heated up provided that gravitational energy is converted into heat. It is widely assumed that thermal ionization can trigger the MRI if the temperature exceeds about 1000 K; however, in CPDs this is not always the case. First of all, the ionization fraction obtained from the Saha equation depends on density, and a gravitationally unstable disk naturally has a high surface density. Second, for a CPD to sustain the MRI, the required ionization fraction is comparatively higher than that of a PPD even if the density is the same.

As is mentioned in Fujii et al. (2014), this is because typical length scale is orders of magnitude smaller in a sub-disk. Thus, the critical temperature is higher than 1000 K. If we adopt \( \eta = 234(T/1K)^{1/2} x_e^{-1} \text{cm}^2 \text{s}^{-1} \) where \( x_e \equiv n_e/n_n \) is ionization degree and \( n_e \) and \( n_n \) are respectively number density of electron and neutral gas (Blaes & Balbus 1994), the condition for sufficient ionization to sustain the MRI at midplane can be given as

\[ \Lambda = \frac{2c_s^2 x_e}{234 \sqrt{T/1K} \Omega_K \beta_0} > 1, \]

where \( \beta_0 \) is the plasma beta at the midplane. From this equation, we can derive the critical ionization degree needed to be MRI-active as \( x_e,\text{crit} = 234 \sqrt{T/1K} \Omega_K \beta_0/(2c_s^2) \).

As an illustration of this, Figure 1 shows the critical ionization degree needed to have the MRI at the midplane for a disk with \( \beta_0 = 10^5 \). The plot is made for the disk temperatures of 100 K, 1000 K, and 10,000 K, respectively. Note that, with the “gas-starved” case of surface density as the lower limit, \( \Sigma = 100 \left(r/20R_J\right)^{-3/4} \text{g cm}^{-2} \) (Canup & Ward 2002; Sasaki et al. 2010), midplane ionization degree at 10 \( R_J \), when \( T = 1000 \) K is assumed, is only about \( 10^{-12} \). One can easily see from Figure 1 that this value is far below the critical value, \( x_e,\text{crit} \), at the respective radius. The temperature needed to obtain \( x_e > x_e,\text{crit} \) for this case is about 2000 K. Obviously, if the surface density is larger, higher temperature is required.

In this paper, we consider the thermal ionization of potassium, sodium, and magnesium. In the high-temperature regime, where most of the metals are already ionized, atomic hydrogen also becomes a dominant source of free electrons. At such temperatures, hydrogen gas is already dissociated, so we employ the following Saha equation for atomic hydrogen as well as for metals to solve for the ionization degree from collisions:

\[
\frac{x_e}{x_K} = 2 \frac{2\pi m_e k_B T_e}{n_n} \left(\frac{2\pi m_e k_B T_e}{h^2}\right)^{3/2} \exp\left(-5.0 \times 10^3/T_e\right),
\]

\[
\frac{x_e}{x_{Na}} = 2 \frac{2\pi m_e k_B T_e}{n_n} \left(\frac{2\pi m_e k_B T_e}{h^2}\right)^{3/2} \exp\left(-6.0 \times 10^4/T_e\right),
\]

\[
\frac{x_e}{x_{Mg}} = 2 \frac{2\pi m_e k_B T_e}{n_n} \left(\frac{2\pi m_e k_B T_e}{h^2}\right)^{3/2} \exp\left(-8.9 \times 10^4/T_e\right),
\]

\[\text{Figure 1. Critical ionization degree needed to sustain the MRI at the midplane at } T = 100 \text{ K, } 1000 \text{ K, and } 10,000 \text{ K for } \beta_0 = 10^5.\]

\[\text{These expressions are approximately correct when the ionization fraction of each species is small, which is not appropriate for K and Na in this case. Since their abundance is small, however, the resulting ionization degree is not strongly affected.}\]
\[
\frac{x_e x_{\text{H}^+}}{x_{\text{H}}} = \frac{2}{n_n} \left( \frac{2\pi m_e k_B T_e}{h^2} \right)^{3/2} \times \exp\left(-1.6 \times 10^5 / T_e\right),
\]
where \(x_K\), etc., represent the abundances of each species, and \(m_e\), \(h\), and \(k_B\) are respectively electron mass, the Planck constant, and the Boltzmann constant. With those equations and charge neutrality, we finally obtain the ionization degree as
\[
x_e = \left( \frac{2}{n_n} \right)^{1/2} \left( \frac{2\pi m_e k_B T_e}{h^2} \right)^{3/4} \times \left[ x_K \exp\left(-5.0 \times 10^4 / T_e\right) \right.
\nonumber
+ x_{\text{Na}} \exp\left(-6.0 \times 10^4 / T_e\right)
\nonumber
+ x_{\text{Mg}} \exp\left(-8.9 \times 10^4 / T_e\right)
\nonumber
+ x_{\text{H}} \exp\left(-1.6 \times 10^5 / T_e\right) \right]^{1/2}. \tag{11}
\]

We adopt solar abundance multiplied by depletion \(\delta\) for metal species, \(x_K = 9.87 \times 10^{-5} \delta\), \(x_{\text{Na}} = 1.60 \times 10^{-6} \delta\), and \(x_{\text{Mg}} = 3.67 \times 10^{-5} \delta\), in our calculation.

If \(\Lambda > 1\) is satisfied at the midplane, we set the viscous parameter due to the MRI turbulence as \(\alpha_{\text{MRI}} = 1040 / \beta_0 + 0.015\) (Okuzumi & Hirose 2011). Thus, we denote the Elsässer number at the midplane as \(\Lambda\) hereafter. What if the ionization degree is not high enough to sustain the MRI and the surface density is smaller than the critical value to be gravitationally unstable? There is always molecular viscosity, but it is negligibly small. Gravitational interaction between the star, planet, and gas of a CPD can be an origin of angular momentum transport (Rivier et al. 2012). Kelvin–Helmholtz-like instabilities between sedimenting dust layers and gas can generate turbulence, which can be roughly estimated as \(\sim 10^{-4} - 10^{-3}\). Moreover, it has been found that in disks with imposed radial temperature gradients, the resulting vertical shear can be a robust source of turbulence via an analog of the Goldreich–Schubert–Fricke instability (Nelson et al. 2013). Rigorously establishing the presence of the vertical-shear instability in CPDs will, however, require us to derive constraints on radiative cooling timescales \(\tau_{\text{crit}}\) similar to the work by Lin & Youdin (2015), who (in the context of PPDs) find the corresponding criterion to scale with the disk thickness—which is favorably large for the comparatively puffed-up CPDs, implying less-restrictive conditions on \(\tau_{\text{crit}}\).

There may be other ways of transporting angular momentum; however, those mechanisms contain uncertainties, and the specific value is not yet obtained. Thus, we treat them via setting a floor value, \(\alpha_{\text{floor}}\), in this work. In summary, we define the viscous parameter as
\[
\alpha = \alpha_{\text{MRI}} + \alpha_{\text{GI}} + \alpha_{\text{floor}} \tag{12}
\]
\[
\alpha_{\text{MRI}} = \begin{cases} 1040 / \beta_0 + 0.015 & (\text{if } \Lambda > 1), \\ 0 & (\text{if } \Lambda \lesssim 1). \end{cases} \tag{13}
\]

We take \(\varepsilon\) and \(\alpha_{\text{floor}}\) as parameters and obtain structures of CPDs based on Section 2.1 above.

3. Resulting Disk Structure

Since the timescale for \(\varepsilon\) to drop is uncertain, we simply develop a disk for each parameter set from scratch until it reaches a quasi-stationary state. First, we calculate assuming only 1% of metals are in gas phase, i.e., \(\delta = 0.01\). Figure 2 is an example of the formation of a CPD with \(\varepsilon = 10^{-2}\) and \(\alpha_{\text{floor}} = 10^{-4}\). Both surface density and temperature increase with time. The uppermost lines in Figure 2 show the values in the steady state except for within \(\sim 2R_h\), where the radial profiles of the surface density and temperature remain non-steady and wiggle about.

In Figure 3, the opacity, Elsässer number, \(\Lambda\), and \(Q\) value at the final stage in Figure 2 are shown. In the quiescent outer disk, the Elsässer number and \(Q\) parameter remain in the stable regime, that is, \(\Lambda < 1\) and \(Q > 2\), and accordingly, \(\alpha\) is determined by \(\alpha_{\text{floor}}\). The Elsässer number occasionally exceeds unity in the inner disk and that prevents the system from settling into a stationary solution. Even if the inner disk remains time-variable, the outer disk achieves a steady state independent of the inner region. In the quiescent outer disk, a bump in the surface-density profile forms because of the increase of opacity due to the transition of the origin from dust sublimation to molecules. The dips in surface density (and temperature) at the inner domain border are related to the boundary conditions. Thus, we do not consider them as a bump.

Figure 4 shows the surface-density profiles (left column) and temperature (right column) of various disk models once the outer disk has reached a steady state. One can see that models with larger \(\varepsilon\) and/or smaller \(\alpha_{\text{floor}}\) generally become more massive and hotter. The top panels of Figure 4 show the surface-density and temperature structure for \(\varepsilon = 10^{-1}\). Because the ionization degree reaches a near-critical value at the inner disk radii, the gas accretion rate fluctuates in time (illustrated by the shaded area in Figure 4), and the disk structure is not fully stationary in this regime. When MRI enabled by thermal ionization is developed, depending on the settings, the disk either ends up with a steady state with smaller surface density or enters the gravito-magneto limit cycle studied by Lubow & Martin (2012, 2013).

The middle panels of Figure 4 show the disk structure for \(\varepsilon = 10^{-2}\). For the case of \(\alpha_{\text{floor}} = 10^{-3}\) (green dotted line), relatively effective gas accretion keeps the surface density smaller than the cases for \(\alpha_{\text{floor}} = 10^{-4}\) (blue solid line) and \(\alpha_{\text{floor}} = 10^{-5}\) (pink dashed line). Since the temperature does not become high enough with \(\alpha_{\text{floor}} = 10^{-3}\), the MRI is not triggered by thermal ionization, and the whole disk settles into a steady state. Compared to the top panels, values are generally slightly smaller.

The bottom panel of Figure 4 illustrates a case where the reduction factor decreases down to \(10^{-3}\). Because the infall flux is already small enough, the surface density does not become that massive, and therefore the temperature cannot be as high as supplying sufficient ionization to sustain the MRI. We remark that the surface-density range of our models is similar to the extended outer disk of Mosqueira & Estrada (2003a), but temperature is much higher in our models. For some parameter sets, a bump can be seen in surface density that is formed due to the change in opacity. A radial pressure bump
cannot be seen in a disk with small $\varepsilon$ and/or large $\alpha_{\mathrm{doo}}$ because the surface density does not pile up sufficiently for the inner disk to transition into the higher temperature regime required for the opacity transition.

Next, for the case when the temperature at which the thermal ionization plays a role for gas accretion is as high as the one for grains to evaporate, we calculate disk structures with $\delta = 1$ corresponding to solar abundance. We show the respective results in Figure 5. For models with $\varepsilon = 10^{-3}$ (not show in the figure), we obtained identical profiles as in the case with $\delta = 0.01$ for all values of $\alpha_{\mathrm{doo}}$. This is also true for the model with $\varepsilon = 10^{-2}$ and $\alpha_{\mathrm{doo}} = 10^{-3}$. For all other models, the range in which we do not obtain steady-state solutions slightly increase because thermal ionization can provide more electrons at the same temperature.

We conclude that the transition of the opacity regime from sublimation of dust to molecules can produce interesting structures in embedded CPDs for a variety of reasonable disk models. In the following section, we discuss the orbital evolution of (proto-)satellites in the disk models derived here.

4. Capture of Satellites in Resonant Orbits

In the context of protoplanetary disks, disk–planet interaction has been studied extensively in the literature. Planets are believed to migrate in the host disk by exchanging their angular momentum with the disk. The idea is also introduced in satellite formation (for instance in Canup & Ward 2002, 2006; Sasaki et al. 2010; Ogihara & Ida 2012). As for PPDs, the migration direction and speed depend on the satellite mass and the disk structure, and the timescale is given as

$$\tau_m = \frac{1}{|b|} \left( \frac{M_s}{M_p} \right)^{1/2} \left( \frac{\Sigma r^2}{M_p} \right)^{1/2} \left( \frac{c_s}{\sqrt{\nu \kappa}} \right)^2 \Omega_{K}^{-1},$$

where $M_s$ and $M_p$ are the mass of the satellite and planet, respectively, and $b$ is a constant that determines the direction and speed of the migration (Paardekooper et al. 2011; Kretke & Lin 2012; Ogihara et al. 2015). In this work, we only consider moons with circular orbits, and we use the formula of Paardekooper et al. (2011) for the migration constant $b$ (see also Ogihara et al. 2015):

$$b = \frac{2}{\gamma} (1 - 2.5 - 1.7q + 0.1p)
+ 1.1F(P_s)G(P_p) \left( \frac{3}{2} - p \right)
+ 0.7(1 - K(P_s)) \left( \frac{3}{2} - p \right)
+ 7.9q(\gamma - 1) \rho F(P_s)F(P_p) \sqrt{G(P_s)G(P_p)}
\times \left( 2 - \frac{1.4}{\gamma} \right) [q - (\gamma - 1)p]
\times \sqrt{(1 - K(P_s))(1 - K(P_p))},$$

where $\gamma$ is the adiabatic constant, $p \equiv -d \ln \Sigma/d \ln r$, and $q \equiv -d \ln T/d \ln r$. The expressions for $F(P)$, $G(P)$, and $K(P)$ are furthermore given as

$$F(P) = \left( 1 + \left( \frac{P}{1.3} \right)^2 \right)^{-1},$$

$$G(P) = \begin{cases} 
\frac{16}{25} \left( \frac{45 \pi}{8} \right)^{3/4} P^{3/2} & P < \frac{8}{45 \pi}, \\
1 - \frac{9}{25} \left( \frac{8}{45 \pi} \right)^{4/3} P^{-8/3} & \frac{8}{45 \pi} \leq P.
\end{cases}$$

$$K(P) = \begin{cases} 
\frac{16}{25} \left( \frac{45 \pi}{38} \right)^{3/4} P^{3/2} & P < \frac{38}{45 \pi}, \\
1 - \frac{9}{25} \left( \frac{28}{45 \pi} \right)^{4/3} P^{-8/3} & \frac{38}{45 \pi} \leq P.
\end{cases}$$

We assume $\gamma = 1.4$ and that the thermal diffusivity is the same as $\nu$ (i.e., $Pr = 1$). Thus, with the dimensionless half-width of the horseshoe region, $\chi_h = 1.1/\gamma^{1/4} \sqrt{M_* r / M_p b}$ (where $h$ is the scale height of the disk),

$$P_p = \frac{2}{3} \sqrt{\frac{\Omega_K r^2 \chi_h^3}{2 \pi \nu}} = \frac{2}{3} P,$$

If $b$ is negative, the satellite migrates toward the planet, and positive $b$ means outward migration. We selected disk models with a discernible bump in surface-density structure, as summarized in Table 2. In the following, we refer to the cases
3. If a satellite migrates and stays at \( b \simeq 0 \), the second satellite migrating from the outer disk approaches the first one that is trapped inside the location of convergent migration. If the migration timescale of the second satellite is longer than the critical timescale, \( t_m^{\text{crit}} \), at the location of the first satellite, the second satellite is captured in a mean motion resonance. As mentioned in Ogihara & Kobayashi (2013), the capture probability for higher-order resonances is very low, and moreover the 2:1 MMR is the outermost among first-order resonances—we hence exclusively focus on this case. The critical timescale for capture into 2:1 MMR of equal-mass satellites is given by (Ogihara & Kobayashi 2013)

\[
t_m^{\text{crit}} = 2.5 \times 10^4 \left( \frac{M_s}{M_{\text{Io}}} \right)^{-4/3} \left( \frac{M_s}{M_{\text{J}}} \right)^{4/3} T_K,
\]

where \( M_{\text{Io}} \) is the mass of Io and \( T_K \) is the orbital period of the satellite. Here, we only consider satellites with equal masses because Galilean satellites have similar masses.

In Model 1, the first satellite is located at about 8.5\text{R}_J after the termination of migration. A satellite in 2:1 resonance with the satellite at 8.5\text{R}_J has an orbit at approximately 14\text{R}_J. The corresponding migration timescale is \( t_m(14\text{R}_J) \simeq 1500 \) years, which is longer than the critical timescale for the first satellite, \( t_m^{\text{crit}}(8.5\text{R}_J) = 97 \) years. This means that the second satellite is captured in the resonance. Similarly, the third satellite can be captured in the 2:1 resonance of the second satellite at 21\text{R}_J because \( t_m(21\text{R}_J) \simeq 2600 \) years > \( t_m^{\text{crit}}(14\text{R}_J) \simeq 200 \) years. In this way, we successfully build up a system in Laplace resonances.

The positions where the first satellite terminate for Models 2–4 are 4.8\text{R}_J, 6.2\text{R}_J, and 3.4\text{R}_J, and the 2:1 resonance orbits of these are 7.6\text{R}_J, 9.8\text{R}_J, and 5.4\text{R}_J, respectively. The orbits in 2:1 resonance with the second satellites are 12\text{R}_J, 16\text{R}_J, and 8.6\text{R}_J, respectively. As summarized in Table 3, the migration timescale of each of these orbits is larger than the critical timescale for capture in the mean motion resonance. Thus, we can also obtain systems in the Laplace resonance with Models 2–4, as well as Model 1. Similarly, we can form those systems in Models 1’, 2’, and 4’. The comparison of orbits of Models 1–4, 1’, 2’, and 4’ with the Galilean moons are given in Figure 8. Note that the orbits of the resonant three moons are located on the same slope of the surface-density profile in all models. One can see that Model 3 has a similar set of orbits with the inner three moons of the Galilean system.

5. Discussion

We summarize the orbits of satellites in our Model 3 along with the mass and orbits of the inner three Galilean moons that are in the Laplace resonance in Table 4. Systems in other models are more compact or spread out compared with the Galilean moons, but most importantly, the moons are in the 4:2:1 MMR. Once they are in this resonance, the orbits are locked, and the moons migrate together as a system; the separations of the bodies adjust accordingly, when the whole system moves radially during the evolution of the CPD.

Figures 4 and 5 show that the disk is quite hot at this stage. At such high disk temperatures, the radial pressure gradient may lead to sub-Keplerian rotation velocities. Actually, as Figure 9 shows, the angular velocity is smaller than the Keplerian value in the outer part of the disk. The angular velocity is calculated as \( \Omega = \Omega_K(1 - \eta)^{1/2} \), where

![Figure 3](image-url)  
Figure 3. Radial profiles of opacity (top), Elsässer number (middle), and Toomre’s Q parameter (bottom panel) for the fiducial case with \( \varepsilon = 10^{-2} \) and \( \alpha_{\text{disc}} = 10^{-7} \).
Takeuchi & Lin (2002). We assumed Keplerian rotation profiles when we derived disk models. However, since Equation (1) is only sensitive to the radial slope of the angular velocity, the assumption is expected to be acceptable.

Hot CPDs are suggested by Keith & Wardle (2014), Zhu (2015), and Szulágyi et al. (2016), but it may be difficult to form icy satellites in such an environment. Since the outer disk is cooler, moons may gain icy materials simply by migrating in from larger radii. Although Figure 6 shows $b$ is positive in $r > 25R_\text{J}$ for Model 4, for instance, bodies about 10 times smaller than Io can migrate all the way from the outer radii because $b$ for them remains negative at all radii. Another possibility is that ice-rich planetesimals are captured when they enter into the CPD. Tanigawa et al. (2014) found that the orbits of sub-Io-sized planetesimals captured in a CPD are highly eccentric. They also found that 10 m or larger planetesimals can be efficiently captured in a CPD, thus those bodies may grow into the size of present moons.

One problem is, however, whether the system can survive over the long-term evolution of the CPD. As mass infall decreases, the temperature of the disk also decreases. When the disk structure that traps the innermost body disappears, the satellite system will start to migrate toward the planet. Moons can survive if the CPD is quickly cleared before they are lost into the planet. A rough estimate of the viscous timescale of the disk is $r^2/\nu \sim 100$ years, which is shorter than the migration timescale of the satellites. However, the actual timescale for the
surface density to become small enough not to affect satellite migration is most likely to be much longer than this estimate. Clearly, this depends on how the disk dissipates and many other unknown factors. In order to obtain a better understanding of how the infall terminates, we need to further study the evolution of PPDs, including both gap formation and gas dissipation. In this work, we adopted a mass infall rate derived from isothermal hydrodynamic simulations, however, Gressel et al. (2013) suggested that taking magnetic field and radiative

Table 2

| Model Parameters for the Cases Considered for Orbital Migration |
|----------------------|----------|----------|
|                      | \( \varepsilon \) | \( \alpha_{\text{floor}} \) |
| Models 1 and 1'      | \( 10^{-1} \)     | \( 10^{-4} \)     |
| Models 2 and 2'      | \( 10^{-1} \)     | \( 10^{-3} \)     |
| Models 3 and 3'      | \( 10^{-2} \)     | \( 10^{-5} \)     |
| Models 4 and 4'      | \( 10^{-2} \)     | \( 10^{-4} \)     |
| Models 5 and 5'      | \( 10^{-3} \)     | \( 10^{-5} \)     |

Figure 5. Same as Figure 4, but for \( \delta = 1 \), that is, undepleted composition. The top panels are for \( \varepsilon = 10^{-1} \), and the bottom panels are for \( \varepsilon = 10^{-2} \). The results for \( \varepsilon = 10^{-3} \) are not shown because they are identical to Figure 4.

Figure 6. Migration coefficient for an Io-sized Moon in the disk model given in Section 2 for Models 1–5. The positions where the migration stops for each model are indicated by crosses.

Figure 7. Migration coefficient for an Io-sized Moon for Models 1'–5'. For comparison, the positions where the first moons stop migration for Models 1, 2, and 4 are marked with crosses.
Comparison of the resonant orbits of the satellites obtained in our models with those of the Galilean moons.

Figure 8.

Table 3
Comparison between Type-I Satellite Migration Timescales, $t_{\text{m}}$, and Critical Timescales, $t_{\text{crit}}^{i}$, for Capture into MMR

| Model   | Migration timescale | Critical timescale |
|---------|---------------------|--------------------|
| 1       | $t_{\text{m}}(14R_J) = 1500$ years | $t_{\text{crit}}^{i}(8.5R_J) = 97$ years |
|         | $t_{\text{m}}(21R_J) = 2600$ years | $t_{\text{crit}}^{i}(14R_J) = 200$ years |
| 2       | $t_{\text{m}}(7.6R_J) = 5000$ years | $t_{\text{crit}}^{i}(4.8R_J) = 41$ years |
|         | $t_{\text{m}}(12R_J) = 7600$ years | $t_{\text{crit}}^{i}(7.6R_J) = 82$ years |
| 3       | $t_{\text{m}}(9.8R_J) = 760$ years | $t_{\text{crit}}^{i}(6.2R_J) = 60$ years |
|         | $t_{\text{m}}(16R_J) = 1300$ years | $t_{\text{crit}}^{i}(9.8R_J) = 120$ years |
| 4       | $t_{\text{m}}(5.4R_J) = 2500$ years | $t_{\text{crit}}^{i}(3.4R_J) = 24$ years |
|         | $t_{\text{m}}(8.6R_J) = 4000$ years | $t_{\text{crit}}^{i}(5.4R_J) = 49$ years |

Figure 9. Angular velocity of the disk for Model 4. Keplerian frequency is also plotted in the dotted line.

6. Summary

We have modeled massive and comparatively hot CPDs by solving the time evolution of surface density with mass infall from the parental PPD. The mass infall flux was determined based on the high-resolution numerical simulation of Tanigawa et al. (2012), where we have also considered the reduction of the flux caused by the dissipation of the PPD at the location of the sub-disk. The temperature profile of the CPD is derived by the balance of viscous heating and radiative cooling, as well as the radial advection. Since the strength of viscosity is uncertain in the absence of MRI, we employed a parameter to determine the minimum value of the viscosity. We considered the MRI when the Elsässer number exceeds unity due to thermal ionization. We furthermore monitored Toomre’s Q parameter in order to consider effective viscosity when the value becomes lower than about 2. When the evolution is governed by $\alpha_{\text{torq}}$, the system settles into a steady state.

In many previous studies, the critical temperature for the onset of the MRI is assumed to be at about 1000 K. As shown in Figure 1, however, we found that this is not the case for massive CPDs. This is for two reasons: (i) the ionization degree needed to sustain the MRI in a CPD is higher than that in a PPD, and (ii) thermal ionization is less effective in higher-density regions. In our models, MRI is turned on by thermal ionization only around $T \sim 2000–3000$ K.

We found that opacity transitions change the radial dependence of the temperature structure, and in particular, a transition near 2000 K makes a bump in surface-density distribution. We estimated whether a moon migrating toward the central planet can be trapped at such a location. In the case of some of the parameter settings that are referred to as Models 1–4, 1′, 2′, and 4′, the surface-density and temperature gradients were sufficiently steep to stop the migration of a moon. Moreover, we have examined the migration timescales of the second and third moons migrating inward and compared them to the critical timescale to be captured in a 2:1 MMR with the inner moon. In all of Models 1–4, 1′, 2′, and 4′, we obtained systems in 4:2:1 mean motion resonance that is known for inner three bodies of the Galilean system. The satellite system obtained in our disk models may or may not survive until the dissipation of the CPD. In order to find out the long-term evolution of these systems, further studies on mass infall from
PPDs and on the origin of angular momentum transport in CPDs are needed.

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References

Armitage, P. J., Livio, M., & Pringle, J. E. 2001, MNRAS, 324, 705
Ayliffe, B. A., & Bate, M. R. 2009, MNRAS, 393, 49
Balbus, S. A., & Hawley, J. F. 1991, ApJ, 376, 214
Bell, K. R., & Lin, D. N. C. 1994, ApJ, 427, 987
Blaes, O. M., & Balbus, S. A. 1994, ApJ, 421, 163
Cannizzo, J. K. 1993, ApJ, 419, 318
Canup, R. M., & Ward, W. R. 2002, AJ, 124, 3404
Canup, R. M., & Ward, W. R. 2006, Natur, 441, 834
Charnoz, S., Salmon, J., & Crida, A. 2010, Natur, 465, 752
Crida, A., & Charnoz, S. 2012, Sci, 338, 1196
Estrada, P. R., & Mosqueira, I. 2006, Icar, 181, 486
Fujii, Y. I., Okuzumi, S., & Inutsuka, S. 2011, ApJ, 743, 53
Fujii, Y. I., Okuzumi, S., Tanigawa, T., & Inutsuka, S. 2014, ApJ, 785, 101
Gressel, O., Nelson, R. P., Turner, N. J., & Ziegler, U. 2013, ApJ, 779, 59
Hayashi, C. 1981, PThPS, 70, 35
Hyodo, R., Charnoz, S., Genda, H., & Ohtsuki, K. 2016, ApJL, 828, L8
Hyodo, R., Charnoz, S., Ohtsuki, K., & Genda, H. 2017, Icarus, 282, 195
Hyodo, R., Ohtsuki, K., & Takeda, T. 2015, ApJ, 799, 40

Keith, S. L., & Wardle, M. 2014, MNRAS, 440, 89
Kimura, S. S., & Tsuribe, T. 2012, PASJ, 64, 116
Klahr, H., & Kley, W. 2006, A&A, 445, 747
Kretke, K. A., & Lin, D. N. C. 2012, ApJ, 755, 74
Lin, M.-K., & Youdin, A. N. 2015, ApJ, 811, 17
Lubow, S. H., & Martin, R. G. 2012, ApJL, 749, L37
Lubow, S. H., & Martin, R. G. 2013, MNRAS, 428, 2668
Lunine, J. I., & Stevenson, D. J. 1982, Icar, 52, 14
Machida, M. N., Kokubo, E., Inutsuka, S., & Matsumoto, T. 2010, MNRAS, 405, 1227
Martin, R. G., & Lubow, S. H. 2011, ApJL, 740, L6
Martin, R. G., & Lubow, S. H. 2013, MNRAS, 432, 1616
Mosqueira, I., & Estrada, P. R. 2003a, Icar, 163, 198
Mosqueira, I., & Estrada, P. R. 2003b, Icar, 163, 232
Nelson, R. P., Gressel, O., & Umurhan, O. M. 2013, MNRAS, 435, 2610
Ogihara, M., & Ida, S. 2012, ApJ, 753, 60
Ogihara, M., & Kobayashi, H. 2013, ApJ, 775, 34
Ogihara, M., Kobayashi, H., Inutsuka, S.-i., & Suzuki, T. K. 2015, A&A, 579, A65
Okuzumi, S., & Hirose, S. 2011, ApJ, 742, 65
Okuzumi, S., & Ormel, C. W. 2013, ApJ, 771, 43
Paardekooper, S.-J., Baruteau, C., & Kley, W. 2011, MNRAS, 410, 293
Perez, S., Dunhill, A., Casassus, S., et al. 2015, ApJL, 811, L5
Rivier, G., Crida, A., Morbidelli, A., & Brouet, Y. 2012, A&A, 548, A116
Sano, T., & Miyama, S. M. 1999, ApJ, 515, 776
Sasaki, T., Stewart, G. R., & Ida, S. 2010, ApJ, 714, 1052
Shakura, N. I., & Sunyaev, R. A. 1973, A&A, 24, 337
Suzuki, T. K., Muto, T., & Inutsuka, S. I. 2010, ApJ, 715, 1289
Szlágyi, J., Masset, F., Lega, E., et al. 2016, MNRAS, 460, 2853
Szlágyi, J., Mayer, L., & Quinn, T. 2017, MNRAS, 464, 3158
Szlágyi, J., Morbidelli, A., Crida, A., & Masset, F. 2014, ApJ, 782, 65
Takahashi, S. Z., Inutsuka, S.-i., & Machida, M. N. 2013, ApJ, 770, 71
Takeuchi, T., & Lin, D. N. C. 2002, ApJ, 581, 1344
Tanigawa, T., Maruta, A., & Machida, M. N. 2014, ApJ, 784, 109
Tanigawa, T., Ohtsuki, K., & Machida, M. N. 2012, ApJ, 747, 47
Tanigawa, T., & Watanabe, S.-i. 2002, ApJ, 580, 506
Turner, N. J., Lee, M. H., & Sano, T. 2014, ApJ, 783, 14
Zhu, Z. 2015, ApJ, 799, 16
Zhu, Z., Hartmann, L., & Gammie, C. 2010, ApJ, 713, 1143