PLANET FORMATION IN SMALL SEPARATION BINARIES: NOT SO SECULARLY EXCITED BY THE COMPANION

Roman R. Rafikov
Department of Astrophysical Sciences, Princeton University, Ivy Lane, Princeton, NJ 08540, USA; rrr@astro.princeton.edu

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ABSTRACT

The existence of planets in binaries with relatively small separations (around 20 AU), such as α Centauri or γ Cephei, poses severe challenges to standard planet formation theories. The problem lies in the vigorous secular excitation of planetesimal eccentricities at separations of several AU, where some of the planets are found, by the massive, eccentric stellar companions. High relative velocities of planetesimals preclude their growth in mutual collisions for a wide range of sizes, from below 1 km up to several hundred km, resulting in a fragmentation barrier to planet formation. Here we show that, for the case of an axisymmetric circumstellar protoplanetary disk, the rapid apsidal precession of planetesimal orbits caused by the disk gravity acts to strongly reduce the direct secular eccentricity excitation by the companion, lowering planetesimal velocities by an order of magnitude or even more at 1 AU. By examining the details of planetesimal dynamics, we demonstrate that this effect eliminates the fragmentation barrier for in situ growth of planetesimals as small as \( \lesssim 10 \) km even at separations as wide as 2.6 AU (the semimajor axis of the giant planet in HD 196885), provided that the circumstellar protoplanetary disk has a small eccentricity and is relatively massive, \( \sim 0.1 \, M_\odot \).

Key words: binaries: close – planets and satellites: dynamical evolution and stability – planets and satellites: formation – protoplanetary disks

1. INTRODUCTION

About 20% of planets detected via stellar radial velocity variations reside in binaries (Desidera & Barbieri 2007). The majority of these systems are wide-separation binaries, with semimajor axis \( a_b \gtrsim 30 \) AU. At the same time, four relatively small separation binaries with \( a_b \approx 20 \) AU (HD 196885, γ Cephei, Gl 86, and HD 41004; Chauvin et al. 2011) are also known to harbor giant planets with projected masses \( M_{pl} \sin i \approx (1.6–4) \, M_J \). In these systems the mass of the secondary star (the binary component other than the star orbited by the planet, which we denote as “primary”) \( M_2 \) is found to be close to 0.4 \( M_\odot \) and the binary eccentricity \( e_b \) is close to 0.4. Also, Dumusque et al. (2012) have recently announced an Earth mass companion to α Centauri B, a member of the binary (or, possibly, a triple) with \( a_b = 17.6 \) AU, \( e_b = 0.52 \), and \( M_b = 1.1 \, M_\odot \). In this system the planet orbits the star at \( \approx 0.04 \) AU separation.

The uniqueness of these systems lies in the fact that planets forming in them are known to provide an extreme challenge to planet formation theories (Zhou et al. 2012). With the exception of α Cen and Gl 86, planets in these binaries reside in rather wide orbits, with planetary semimajor axes of \( a_{pl} \approx 1.6–2.6 \) AU. In situ formation of these gas giants is expected to proceed through the continuous agglomeration of planetesimals, starting from small objects (easily \( \lesssim 1 \) km). However, the gravity of the eccentric stellar companion inevitably leads to rapid secular evolution (Heppenheimer 1978), driving planetesimal eccentricities far above the level at which bodies are not destroyed in mutual collisions (Thébault et al. 2008). This problem, which is often called the collisional or fragmentation barrier, is especially severe for small (1–10 km) planetesimals for which the ratio of binding to kinetic energy is small, and far from the primary, where secular forcing by the companion is strongest.

Marzari & Scholl (2000) suggested that a combination of secular forcing by the companion and gas drag acting on small (1–10 km) planetesimals leads to an apsidal alignment of their orbits, resulting in smaller collision velocities and enabling growth. However, Thébault et al. (2006, 2008) demonstrated that size-dependence of apsidal alignment acts to break orbital phasing between objects of different sizes, resulting in high-velocity collisions and reinforcing the fragmentation barrier.

Most studies of planetesimal growth in binaries include the effect of the protoplanetary disk on planetesimal dynamics only via gas drag (Thébault et al. 2004, 2006, 2008, 2009; Paardekooper et al. 2008; Paardekooper & Leinhardt 2010), without accounting for the gravitational effect of the disk. Fragner et al. (2011), Zhao et al. (2012), and Batygin et al. (2011) have considered disk gravity in the context of planet formation and evolution in systems with misaligned stellar companions (Fragner et al. 2011 also briefly mention the aligned case). This setup is probably irrelevant for small separation (tens of AU) binaries, which we expect to be coplanar with circumstellar disks (Hale 1994). On the other hand, Kley & Nelson (2007) and Marzari et al. (2009) demonstrated that perturbation due to the coplanar companion generically makes the gas disk eccentric. Non-axisymmetric disk potential then excites high planetesimal eccentricities, in addition to direct excitation by the companion, further inhibiting planetesimal growth.

In this Letter we explore planetesimal dynamics in small separation binaries assuming an axisymmetric gas disk. A low-eccentricity disk may be possible if it is sufficiently compact (which reduces tidal forcing by the companion) and massive (self-gravity lowers disk eccentricity; Marzari et al. 2009), which is consistent with the findings of this work; see Section 4. In this setup we show that the apsidal precession of planetesimal orbits induced by disk gravity dominates secular evolution at separations of several AU and the relative velocities at which bodies collide are reduced, sometimes by more than an order of magnitude. In massive axisymmetric disks this may naturally eliminate the fragmentation barrier issue for the in situ formation of giant planets in systems such as γ Cephei.

2. SECULAR EVOLUTION

We consider planetesimal motion as Keplerian motion around the primary perturbed by the gravity of the companion, which
moves on a larger, eccentric orbit, and of the disk. The mass of the primary is $M_p$, and we define $\mu \equiv M_\ast/(M_p + M_\ast)$. We assume the eccentricity of the stellar binary $e_b$ to be small and the planetesimal orbit to be coplanar with the binary. Planetesimals are immersed in a massive, axisymmetric gaseous disk, characterized by surface density $\Sigma(r)$, as specified in Section 2.1.

Assuming $e \ll 1$, the secular (averaged over the planetesimal and binary orbital motion) disturbing function for a planetesimal with a semimajor axis $a$ and eccentricity vector $\mathbf{e} = (k, h) = (e \cos \sigma, e \sin \sigma)$ (with the apsidal angle $\sigma$ counted from the binary apsidal line, which is assumed to be fixed\(^1\)) is (Murray & Dermott 1999)

$$R = na^2 \times \left[ \frac{1}{2} (A + \dot{\sigma}_d) (h^2 + k^2) - Bk \right], \quad (1)$$

where

$$A = \frac{3}{4} \mu \frac{n_0^2}{n} \left( 1 + \frac{3}{2} e_b^2 \right), \quad (2)$$

$$B = \frac{15}{16} e_b n_0 \frac{a}{n} \frac{a_b}{a} \left( 1 + \frac{5}{2} e_b^2 \right), \quad (3)$$

and

$$\dot{\sigma}_d = -\frac{1}{2n} \left( \frac{2}{r} \frac{\partial U_d}{\partial r} + \frac{\partial^2 U_d}{\partial r^2} \right) \bigg|_{r=a} \quad (4)$$

is the precession frequency of the planetesimal orbit due to the disk potential $U_d$. Here $n_0 = (G(M_p + M_\ast)/a_{\text{tr}}^3)^{1/2}$ and $n = (GM_p/a^3)^{1/2}$ are the mean motions of the binary and planetesimal, respectively. The contribution to $R$ proportional to $\dot{\sigma}_d$ arises from the expansion of $U_d$ along the eccentric planetesimal orbit and averaging over the mean longitude. The axisymmetry of the disk is important since the gravity of an eccentric disk would contribute to planetesimal eccentricity excitation, which is not included here. Evolution equations for $k$ and $h$ are written using $dh/dt = (na^2)^{-1} \partial R/\partial k$, $dk/dt = -(na^2)^{-1} \partial R/\partial h$ as

$$\frac{dh}{dt} = (A + \dot{\sigma}_d)k - B, \quad \frac{dk}{dt} = -(A + \dot{\sigma}_d)h. \quad (5)$$

These equations agree with Marzari & Scholl (2000) if $\dot{\sigma}_d = 0$.

We write the solution for $\mathbf{e} = (e_{\text{free}}(t) + e_{\text{forced}}(t))$, where

$$\begin{align*}
\begin{cases}
k_{\text{free}}(t) \\
h_{\text{free}}(t)
\end{cases} &= e_{\text{free}} \begin{cases}
\cos \left[ (A + \dot{\sigma}_d) t + \sigma_0 \right] \\
\sin \left[ (A + \dot{\sigma}_d) t + \sigma_0 \right]
\end{cases}, \\
\sigma_0 &= \text{a constant, and}
\end{align*}$$

$$\begin{align*}
\begin{cases}
k_{\text{forced}}(t) \\
h_{\text{forced}}(t)
\end{cases} &= e_{\text{forced}} \begin{cases}
1 \\
0
\end{cases}, \\
e_{\text{forced}} &= B/(A + \dot{\sigma}_d). \quad (6)
\end{align*}$$

Thus, the free eccentricity vector $e_{\text{free}}$ rotates at a rate of $A + \dot{\sigma}_d$ around the endpoint of the fixed forced eccentricity vector $e_{\text{forced}}$. Setting $\dot{\sigma}_d = 0$ we reproduce the solution obtained by Heppenheimer (1978).

Planetesimals starting on circular orbits have $e_{\text{free}} = e_{\text{forced}}$ so that their eccentricity oscillates with an amplitude of $2e_{\text{forced}}$ and a period of $2\pi/(A + \dot{\sigma}_d)$.

\(^1\) Precession period of the secondary orbit exceeds important timescales of the problem.

### 2.1. Disk Model

We model the disk as a constant $M$ disk extending to the outer truncation radius $r_t$. Numerical simulations of accretion disks in binaries suggest that $r_t \sim (0.2\text{–}0.4) a_b$ (Zhou et al. 2012), depending on $e_b$ and $\mu$. In our study we take $r_t = 0.25 a_b$.

The assumption of constant $M$ is a necessary simplification, which ignores the details of the disk structure at $r \sim r_t$. Assuming that the disk viscosity $\nu$ is well characterized by the radius-independent effective $\alpha$-parameter (Shakura & Sunyaev 1973), we write

$$\Sigma(r) = \frac{M}{3\pi \nu} = \frac{\Omega M}{3\pi \alpha c_s^2}. \quad (8)$$

If disk temperature scales as $T \propto r^{-q}$, then

$$\Sigma(r) = \Sigma_0 \left( \frac{r}{r_0} \right)^p, \quad p = \frac{3}{2} - q. \quad (9)$$

where $r_0$ is some fiducial radius and $\Sigma_0 = \Sigma(r_0)$.

Models of protoplanetary disks typically find $q$ to be close to $1/2$, so that $p \approx 1$. In particular, the passive disk model of Chiang & Goldreich (1997) has $q = 3/7$, so that $p = 1 - 1/14$. For simplicity, in our calculations we take $p = 1$, which corresponds to a classical Mestel disk (Mestel 1963) if $r_t \to \infty$.

Assuming that the surface density has profile (9) with $p = 1$ all the way to $r_t$, we can express $\Sigma$ via the total disk mass $M_d$ enclosed within $r_t$ as

$$\Sigma(r) \approx \frac{M_d}{2\pi r_t^2} \approx 2800 \text{ g cm}^{-2} \frac{M_d}{10^{-3} M_\odot} r_t^{-1} \approx 25 \text{ g cm}^{-2} r_t^{-1}, \quad (10)$$

where $r_{t,5} \equiv r_t/(5 \text{ AU})$, and $r_1 \equiv r/(1 \text{ AU})$.

### 2.2. Precession Due to the Disk

A disk with the density profile (9) with $p = 1$ extending to infinity is known to have a constant circular velocity (Mestel 1963)

$$v_c = \left( \frac{r}{r_0} \right)^{1/2} = (2\pi G \Sigma_0 r_0)^{1/2}. \quad (11)$$

Expressing $\partial U_d/\partial r$ from this relation and plugging it in Equation (4), we find $\dot{\sigma}_d = -\pi G \Sigma(r)/(nr)$, where from now on we use $r$ instead of $a$. Even though in our problem disk is truncated at $r_t$, this expression should still give a reasonable estimate of the precession rate $\dot{\sigma}_d$ due to disk gravity for $r \lesssim r_t$.

Using Equation (10), we find

$$\dot{\sigma}_d \approx -\frac{G M_d}{n r_t r^2} = -\frac{n M_d}{M_p} \frac{r}{r_t}. \quad (12)$$

Note that $\dot{\sigma}_d$ varies rather weakly with $r$, as $r^{-1/2}$, which is consistent with Batygin et al. (2011).

### 2.3. Planetesimal Eccentricities

To assess the role of a disk-driven precession on secular evolution, we evaluate

$$\frac{|\dot{\sigma}_d|}{A} \approx 4 \frac{M_d}{3 M_\ast} \frac{a_b^3}{r_t r^2} \approx 20 \frac{M_d}{M_\ast} \frac{a_b^3}{r_t r^2} \frac{1}{r_{t,5}^2}, \quad (13)$$
where $a_{b,20} \equiv a_b/(20 \, \text{AU})$. In making this estimate we neglected the quadratic term in $e_b$ in Equation (2). It is then obvious that the dynamics of planetesimals at several AU are strongly affected by the disk-driven precession: $|\dot{\sigma}_d|$ exceeds $A$ for

$$r \lesssim r_{\text{c}} \approx 4.6 \, \text{AU} \left( \frac{M_d/M_*}{a_{b,20}} \right)^{1/2},$$

(14)
i.e., essentially for any $r \lesssim r_{\text{c}}$ even for the disk mass as small as $\sim 10^{-2} \, M_\odot$. Thus, when studying planet formation at 2–3 AU we can neglect planetesimal precession due to the secondary compared to the disk-driven precession, i.e., neglect $A$ compared to $|\dot{\sigma}_d|$ in Equation (7) and other formulae.

Equation (7) then predicts that the amplitude of eccentricity oscillations is

$$e_{\text{disk}}(r) = \frac{2B}{|\dot{\sigma}_d|} \approx \frac{15}{8} e_{b} \frac{M_d}{M_*} \frac{r^3}{a_b^2} r_i,$$

(15)

$$\approx 3 \times 10^{-3} \frac{e_{b}}{0.5 M_d/M_*} \frac{a_{b}}{a_{b,20}} r_i^3,$$

(16)

where we again neglected the $e_b^2$ term in Equation (3). This is to be compared with

$$e_{\text{h/disk}}(r) = \frac{5}{2} e_{b} \frac{r}{a_b} \approx 6.3 \times 10^{-2} \frac{e_{b}}{0.5 a_{b,20}} r_i,$$

(17)

which one finds by neglecting the disk-driven precession, i.e., dropping $\dot{\sigma}_d$ in Equation (7). It is obvious that neglecting the disk-driven precession leads to an overestimate of the planetesimal eccentricity at ~AU separations by more than an order of magnitude. This has important consequences for planetesimal growth as we discuss further.

2.4. Gas Drag

Equation (15) accounts for the presence of the disk only through the precession caused by its gravity. However, for small planetesimals gas drag is also important. Assuming quadratic gas drag force in the form $F \approx -\rho v_\text{orb}^2/\rho d$ (here $d$ and $\rho$ are the object’s radius and bulk density, $\rho_d \approx \Sigma/h$ is the gas density, and $h$ is the disk scale height), we account for its effect on planetesimal dynamics by adding the terms $-D(h,k)(h^2+k^2)^{1/2}$ to the first and second equations of Equation (5), respectively (Marzari & Scholl 2000). Here

$$D = n \frac{\rho_d}{\rho_d h} = \frac{\Sigma}{\rho_d h} r_i.$$  

(18)

For small planetesimal sizes (to be specified later by Equation (20)) in the gas drag-dominated regime, drag force balances the eccentricity excitation due to the secondary, i.e., the $B$ term in the first Equation (5). This results in the following estimate for the gas drag-mediated planetesimal eccentricity:

$$e_{\text{gas}} \approx \left( \frac{B}{D} \right)^{1/2} = \left( \frac{e_{b} M_d h \rho_d}{M_p r \Sigma} \right)^{1/2} \frac{r_i}{a_b^2},$$

(19)

This expression agrees with Paardekooper et al. (2008) and predicts that $e_{\text{gas}} \propto d^{1/2}$. As a result, for small bodies one finds $e_{\text{gas}} < e_{\text{disk}}$.

The transition between the drag-dominated behavior (Equation 19) and the drag-free eccentricity scaling analogous to (15) is given by $d_{\text{gas}} = \sqrt{e_{\text{gas}}^2 + e_{\text{disk}}^2}$. Figure 1. Characteristic planetesimal sizes vs. radius for two disk masses $M_d$: (a) $0.01 \, M_\odot$ and (b) $0.1 \, M_\odot$. We display $d_{\text{coll}}$ (solid, Equation (25)), $d_{\text{gas}}$ (dotted, Equation (20)), $d_{\text{disk}}$ (short-dashed, Equation (26)), and planetesimal radius $d_i = 10 \, \text{km}$ below which we consider objects as strength-dominated (long-dashed). The two latter sizes do not depend on $M_d$. Calculations are done for $e_{b} = 0.4$, $M_\odot = 0.4 \, M_\odot$, $a_b = 20 \, \text{AU}$, $M_p = M_\odot$ (typical for small separation binaries; Chauvin et al. 2011), and $r_i = 5 \, \text{AU}$, and $r_i/h = 30$. Planetesimals in the shaded region (“danger zone”) get destroyed according to criterion (22) precluding their growth at corresponding separations. The accretion friendly zone is to the left of the vertical dotted line in each plot; it is wider for higher $M_d$ and extends to $\approx 2.5 \, \text{AU}$ for $M_d = 0.1 \, M_\odot$.

(Equation 15) occurs at the planetesimal size $d_{\text{gas}}$, where these two equations yield the same eccentricity:

$$d_{\text{gas}} \approx \frac{4n}{B} \frac{\rho_d}{\rho} \frac{\Sigma}{r_i} \approx \frac{15}{8\pi} \frac{r}{e_b} \frac{M_d}{h} \frac{M_p r}{\rho_d} ,$$

(20)

$$\approx 1 \, \text{km} \frac{e_b}{0.1} \frac{r_i}{h} \frac{M_p}{M_*} \frac{1}{\rho_d} \frac{\Sigma}{r_i} a_{b,20},$$

(21)

where $\rho_d \approx \rho/(3 \, \text{g cm}^{-3})$, $M_{p,b} \equiv M_p/M_{\odot}$, and we have used Equation (10). Using Equation (17) instead of Equation (15) in this estimate would yield much larger $d_{\text{gas}}$, in agreement with previous studies neglecting the disk-driven precession (e.g., Thébault et al. 2006, 2008).

Planetesimal eccentricity behaves according to formula (19) for $d \lesssim d_{\text{gas}}$, and switches to the drag-free regime (Equation 15) for $d \gtrsim d_{\text{gas}}$; see Figure 1. The dependence of $d_{\text{gas}}$ on disk mass $d_{\text{gas}} \propto M_d^{-1}$—is somewhat counterintuitive since higher gas density results in stronger drag, making it more important for larger bodies. However, the gas-free planetesimal eccentricity (Equation 15) is itself a function of $M_d$ and decreases faster with increasing $M_d$ than does $e_{\text{gas}}$, explaining the nontrivial $d_{\text{gas}}(M_d)$ dependence.

3. IMPLICATIONS FOR PLANET FORMATION

Planetesimals grow as long as their encounter velocity $v_\text{coll}$ (measured at infinity) is such that collisions do not result in a net loss of mass. The conditions for this depend, in particular, on planetesimal size and on whether planetesimals are strength- or gravity-dominated. Using the results of Leinhardt & Stewart
that erosion in equal-mass planetesimal collisions is avoided for

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dominate over the gravitational energy, which is expected to

d from 20 AU to 30 AU reduces (2012) for equal-mass strengthless bodies, we estimate the

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(2012) for equal-mass strengthless bodies, we estimate the condition for planetesimal growth to be

\[ v_{\text{coll}} \lesssim 2v_{\text{esc}}, \]  

(22)

where the escape speed from the surface of an object of radius \( d \) and bulk density \( \rho \) is

\[ v_{\text{esc}} = \left( \frac{8\pi}{3} G \rho \right)^{1/2} d \approx 1.3 \text{ m s}^{-1} \rho^{1/3} d \]  

(23)

(here \( d_1 \equiv d/(1 \text{ km}) \)). It becomes harder to break planetesimals when they are small enough for their internal strength to dominate over the gravitational energy, which is expected to happen for \( d \lesssim d_1 \sim 10 \text{ km} \) (Holsapple 1994).

Planetesimal collisions occur at a velocity of the order of

\[ v_{\text{coll}}(r) \approx c_{\text{disk}}v_K, \]  

where \( v_K = nr \) is the Keplerian speed; we neglect a possible apsidal alignment of planetesimal orbits (Marzari & Scholl 2000) since it is effective only when gas drag is important, i.e., only for small objects \( \lesssim d_{\text{gas}} \). Using expression (15) we find

\[ v_{\text{coll}}(r) \approx 90 \text{ m s}^{-1} \frac{e_b}{0.5} \frac{M_{p,1}}{M_s} \frac{r_{1.5}}{a_{b,20}^4} \frac{1/2}{r_1^{5/2}}. \]  

(24)

Substituting Equations (23) and (24) into condition (22) we find that erosion in equal-mass planetesimal collisions is avoided for bodies with \( d \gtrsim d_{\text{coll}} \), where

\[ d_{\text{coll}} \approx 35 \text{ km} \frac{e_b}{0.5} \frac{M_{p,1}}{M_s} \left( \frac{M_{p,1}}{\rho_3} \right)^{1/2} \frac{r_{1.5}}{a_{b,20}^4} \frac{1/2}{r_1^{5/2}}. \]  

(25)

Thus, for the fiducial binary parameters adopted here and for \( M_d \sim 10^{-2} M_\odot \), only planetesimals larger than \( \approx 35 \text{ km} \) would be able to grow at 1 AU. At 2 AU—the semimajor axis of \( \gamma \) Cephei AB—only bodies larger than 200 km in radius would be able to survive in equal-mass collisions.

However, in the absence of a disk (i.e., without both disk gravity and gas drag) the problem is much worse: evaluating collisional velocity as \( v_{\text{coll}} = c_{\text{disk}}v_K \) using Equation (17) and applying condition (22), one finds that only planetesimals larger than

\[ d_{\text{coll}}^{\text{disk}} \approx 700 \text{ km} \frac{e_b}{0.5} \left( \frac{M_{p,1}}{\rho_3} \right)^{1/2} \frac{1/2}{a_{b,20}^4} r_1^{5/2}. \]  

(26)

are able to survive in equal-mass collisions in the absence of disk-induced precession (and drag-induced phasing). Clearly, the collisional barrier appears to be far more severe if one disregards the effects of the disk on the secular evolution of planetesimals. We compare the behavior of \( d_{\text{gas}}, d_{\text{coll}}, \) and \( d_{\text{coll}}^{\text{disk}} \) as a function of \( r \) in Figure 1.

Note that \( d_{\text{coll}} \) is very sensitive to the binary semimajor axis \( a_b \), unlike \( d_{\text{coll}}^{\text{disk}} \); see Equations (25) and (26): increasing \( a_b \) from 20 AU to 30 AU reduces \( d_{\text{coll}} \) by a factor of five. The size of the region where disk-driven precession dominates the secular evolution also expands rapidly with increasing \( a_b \); see Equation (14).

4. PLANETESIMAL ACCRETION IS POSSIBLE IN MASSIVE AND AXISYMMETRIC GAS DISKS

We now propose a solution to the problem of planetesimal accumulation in binaries, which was raised in Section 1.

\[ \frac{M_d}{M_\star} \gtrsim \frac{e_b}{0.5} \left( \frac{M_{p,1}}{\rho_3} \right)^{1/2} \frac{r_1/\alpha_b}{0.25} \frac{a_2^{5/2}}{a_{b,20}^{3/2}}. \]  

(27)

In Figure 2 we illustrate this constraint as a function of \( a_b \), for different values of \( \alpha_b \). It is clear that in very small separation binaries with \( a_b = 10 \text{ AU} \), growing planets even at 1 AU require a massive disk, \( M_d \approx 0.2 M_\star \).

\[ \text{Figure 2. Plot of } \alpha_b - M_d/M_\star \text{ phase space illustrating conditions under which planets can form in binaries, assuming } d_1 = 10 \text{ km}. \]
In giant planet-hosting binaries, companions typically have mass $M_p \approx 0.4 M_\odot$ (Chauvin et al. 2011), meaning that our scenario of planet formation at 2 AU needs $M_d \approx 0.1 M_\odot$. One might worry that such massive disks would be prone to gravitational instability. With the density profile (10), we estimate the Toomre $Q \equiv n c_s/(\pi G \Sigma)$ ($c_s$ is the sound speed) as

$$Q \approx 2 \frac{M_p h}{M_d} \frac{r}{r} \approx 3 \left( \frac{0.1}{M_d/M_p} \right) \frac{(30 \, \frac{r_s}{r}) \, r_s}{r}. \quad (28)$$

Thus, even for $M_d = 0.1 M_p \approx 0.1 M_\odot$, the disk is at most marginally unstable at 2 AU. But, even if it were unstable, the surface density and optical depth at this distance would be so high that the cooling time would far exceed the local dynamical time, making planet formation by direct disk fragmentation impossible (Gammie 2001; Rafikov 2005). Instead, the disk would slowly evolve under the action of gravitoturbulence (Rafikov 2009).

Disk masses $\lesssim 0.1 M_\odot$ may seem high in light of existing observational evidence for low-mass disks in small separation binaries (Harris et al. 2012). Thus, massive disks in close binaries must be rather rare and/or transient objects.

On the other hand, high $M_d$ greatly simplifies planet formation. First, self-gravity of a massive disk acts to reduce disk eccentricity (Marzari et al. 2009), thereby removing an additional source of planetesimal excitation (Kley & Nelson 2007). Note that calculations presented here explicitly assume an axisymmetric disk, i.e., very small disk eccentricity. Second, known giant planets in small separation binaries are quite massive ($\gtrsim M_J$; Chauvin et al. 2011), thus requiring a large gas reservoir (i.e., high $M_d$) for their assembly. Third, a higher disk density means a larger isolation mass (Lissauer 1993) possibly making it high enough to trigger core accretion at 2 AU without going through the long-lasting stage of giant impacts (Chambers 2004). Fourth, a higher $\Sigma$ likely implies a larger dead zone (Gammie 1996) in the disk, providing quiet conditions for planetesimal formation and growth, and resulting in smaller viscosity, which possibly means a longer disk lifetime. Finally, the timescale on which planets form also goes down as $M_d$ increases.

Another potential solution to the collisional barrier problem is the direct formation of large planetesimals by, e.g., streaming and/or gravitational instabilities (Johansen et al. 2007; Thébault 2011). Large $M_d$ and small $d_{\text{coll}}$ are helpful for this mechanism as well since, to overcome the fragmentation barrier in a massive disk, such instabilities only need to produce bodies with sizes of tens of km rather than $d \sim 10^3$ km.

Thus, even though massive disks may not be very common in close binaries, the efficiency of giant planet formation in them can be much higher than in their low-mass analogs due to the aforementioned factors. This may result in a reasonable rate of production of systems such as γ Cephei.

5. DISCUSSION

There are other factors that may affect our conclusions. First, our growth condition (22) is likely too stringent. Previously, using a more refined fragmentation criterion, Thébault (2011) found that planetesimal growth is possible in HD 196885 even in the absence of disk-driven precession at 2.6 AU as long as the planetesimal size exceeds 250 km. However, according to our formula (26) with the same assumptions ($M_p = 1.3 M_\odot$, $a_p = 21$ AU, $e_p = 0.42$, and $\sigma_d = 0$), growth is possible only for $d \gtrsim 10^3$ km. Also, growth may proceed mainly via unequal-mass collisions, which more frequently result in mergers (Thébault 2011). Thus, our fragmentation criterion likely overestimates the planetesimal size above which objects grow efficiently. It might be easier to overcome the fragmentation barrier with a more realistic growth condition than our simple criterion (22), resulting in lower $M_d$.

On the other hand, there are also factors complicating planetesimal growth. Most importantly, even very low eccentricity of the gaseous disk induced by the companion may very adversely affect growth of small bodies. Some non-zero gas eccentricity has so far always been observed in numerical studies of disks in binaries (Kley & Nelson 2007; Paardekooper et al. 2008; Marzari et al. 2009; Zsom et al. 2011; Müller & Kley 2012). Also, the gas-drug-induced inspiral of planetesimals may deplete the disk of some solids. The relative importance of these factors for planet formation in binaries will be assessed in the future.

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