Fragmentation of Kozai–Lidov Disks

Wen Fu¹,²,³, Stephen H. Lubow⁴, and Rebecca G. Martin¹

¹Department of Physics and Astronomy, University of Nevada, Las Vegas, Las Vegas, NV 89154, USA; wr5@rice.edu
²Department of Physics and Astronomy, Rice University, Houston, TX 77005, USA
³Los Alamos National Laboratory, Los Alamos, NM 87545, USA
⁴Space Telescope Science Institute, Baltimore, MD 21218, USA

Received 2016 September 28; revised 2016 December 5; accepted 2016 December 22; published 2017 January 30

Abstract

We analyze the gravitational instability (GI) of a locally isothermal inclined disk around one component of a binary system. Such a disk can undergo global Kozai–Lidov (KL) cycles if the initial disk tilt is above the critical KL angle (of about 40°). During these cycles, an initially circular disk exchanges its inclination for eccentricity, and vice versa. Self-gravity may suppress the cycles under some circumstances. However, with hydrodynamic simulations that include self-gravity, we show that for a sufficiently high initial disk tilts and for certain disk masses, disks can undergo KL oscillations and fragment due to GI, even when the Toomre Q value for an equivalent undisturbed disk is well within the stable regime (Q > 2). We suggest that KL triggered disk fragmentation provides a mechanism for the efficient formation of giant planets in binary systems and may enhance the fragmentation of disks in massive black hole binaries.

Key words: accretion, accretion disks – binaries: general – hydrodynamics – planets and satellites: formation – quasars: general

1. Introduction

The stability of a self-gravitating gaseous disk is typically characterized by the Toomre parameter, \( Q = c_s \kappa / \pi G \Sigma \) (Toomre 1964), where \( c_s \) is the sound speed, \( \kappa \) is the epicyclic frequency, \( G \) is the gravitational constant, and \( \Sigma \) is the disk surface density. The disk becomes unstable to gravitational instability (GI) when \( Q \) is sufficiently small. This dynamical instability manifests itself as multi-armed spiral waves. The disk torques and shocks produced by these spiral structures redistribute mass and angular momentum, and ultimately help to stabilize the disk. With efficient radiative cooling or mass accretion onto the disk, GI can be maintained and leads to the collapse of regions on the spiral arms into gravitationally bound clumps. The exact condition for the development of disk GI into disk fragmentation is still an active area of research. Nevertheless, it is generally agreed that a necessary condition for GI against non-axisymmetric perturbations is \( Q \lesssim 1.5 \) (Papaloizou & Savonije 1991; Nelson et al. 1998; Mayer et al. 2004). For non-isothermal disks, the \( Q \) criterion is not sufficient for fragmentation because fragmentation also depends on the details of the disk thermodynamics (Gammie 2001; Rice et al. 2005; Lodato & Clarke 2011; Paardekooper et al. 2011; Kratter & Lodato 2016).

Disk GI has long been suggested as an alternative theory for giant planet formation (Boss 1997). Compared with the more standard core-accretion theory, it has two advantages. First, massive planets form within a reasonable amount of time and second, it has the ability to form planets at large disk radii (see the review by Helld et al. 2014). However, this mechanism does not work close to the star where the disk cannot cool efficiently (Rafikov 2005) and the possibility of forming planets preferentially in the outer parts of the disk via disk GI has difficulty explaining the majority of the known exoplanet population (Rice et al. 2015).

There has been a plethora of numerical investigations (either with smoothed particle hydrodynamics (SPH) codes or grid-based finite-difference/finite-volume code) devoted to studying various aspects of forming giant planets via disk GI, such as radiative transfer, effects of disk metallicity and chemistry, effects of numerical resolution and clump evolution (e.g., Boss 2009; Galvagni et al. 2012; Rogers & Wadsley 2012; Nayakshin & Cha 2013; Vorobyov et al. 2013; Evans et al. 2015; Meru 2015; Stamatellos 2015; Tsukamoto et al. 2015). Most of these works assume that circumstellar disks form and exist in isolation. However, 40% to 50% of observed exoplanets are estimated to be in binary star systems (Horch et al. 2014).

Most models for GI involve a circular disk that surround a single star. However, perturbations to the disk, such as those due to a binary companion, can potentially aid in producing fragmentation. Some studies have investigated the possibility of forming giant planets by disk GI in binary systems, but they have given conflicting conclusions regarding the role played by the binary companion in the disk GI development (Nelson 2000; Mayer et al. 2005; Boss 2006). These papers modeled disks that are coplanar with the binary orbital plane. While disk coplanarity might be a good assumption under some circumstances, observational evidence suggests that it is not true for all binary systems. Large mutual inclinations (greater than 60°) have been observed between the circumstellar disks around young binary system components (e.g., Jensen & Akeson 2014; Williams et al. 2014). The binary orbital planes in these systems are unknown but at least one of the disks in each system is possibly significantly inclined (>45°) with respect to the binary orbit. Thus, it is important to understand disk GI and possible giant planet formation in inclined circumstellar disks.

When the inclination of a tilted disk is greater than the critical KL angle (of about 40°), the disk can undergo global Kozai–Lidov (KL) oscillations (Martin et al. 2014; Fu et al. 2015a). The oscillations periodically exchange the disk eccentricity and inclination, as occurs in the ballistic particle case (Kozai 1962; Lidov 1962). In Fu et al. (2015b) we took a
step further and included the effects of disk self-gravity. We found that the disk KL mechanism may be suppressed by disk self-gravity. However, the suppression “window” in terms of the disk mass and the disk inclination is quite narrow. The disk KL mechanism may still operate even when the disk is nearly gravitationally unstable, if the inclination is large.

For disk tilts greater than about 60°, the KL effect attempts to induce a large disk eccentricity that cannot be fully accommodated by a smooth continuous disk. Instead, the disk undergoes strong shocks during its KL oscillations. Such shocks could have important consequences for disk fragmentation.

In this Letter, we report on fragmentation triggered by the KL mechanism in the outer disk regions. In the current analysis, we assume that the disk is locally isothermal, as is more likely the case in outer disk regions, and ignore the complexities of non-isothermal effects. Our emphasis is on the dynamical consequences of shocks driven in KL disks that provide much stronger disk perturbations than have been previously investigated (e.g., Boss 2006). We describe our three-dimensional hydrodynamic simulation setup and main results in Section 2. We compare our study with earlier work on coplanar disks and discuss the implications of our results in Section 3.

2. Numerical Simulation

We carry out three-dimensional SPH simulations of a fluid disk that orbits one member of an equal mass binary system. This binary system is initially on a circular orbit with separation $a_b$. The total mass of the binary is $M = M_c + M_p$, where $M_c$ is the mass of the central star and $M_p$ is the mass of the perturber star. The initial disk mass is $M_d = 0.035 M_*$, or 7% of the central stellar mass, $M_c$. In our simulations, both the central star and the perturber star can feel the gravitational force from the disk and this affects the binary orbit. However, this effect is fairly small and the binary orbit remains almost circular throughout the simulation. Initially the orbital plane of the disk is inclined to the binary orbital plane by $i = 70°$ ($i = 0$ would be a coplanar disk). We use a locally isothermal equation of state and an explicit accretion disk viscosity. The sound speed of the disk is $c_s \propto r^{-3/4}$ and the initial surface density distribution is $\Sigma \propto r^{-\gamma}$. These are chosen such that both $\alpha$ (Shakura & Sunyaev 1973) and the smoothing length $\langle h \rangle/H$ are constant over the disk radius, $r$ (Lodato & Pringle 2007). We use $\alpha = 0.01$ in this study, and the disk initially extends from radius $r_{in} = 0.025 a_b$ to $r_{out} = 0.25 a_b$. The initial circular velocity is corrected for the effects of disk self-gravity. The vertical gas density distribution is taken to be that for hydrostatic balance of a non-self-gravitating disk. Although this initial state is somewhat out of vertical force balance due to self-gravity, this imbalance is unlikely to be responsible for the strong effects we find to be associated with large-scale shocks (see also Backus & Quinn 2016).

We treat the stars as sink particles with a boundary condition such that whenever particles move into the accretion radius they are removed from the simulation while their mass and momentum are deposited onto the sink. There are $1 \times 10^6$ SPH particles at the beginning of the simulation. The disk aspect ratio at the inner disk edge is $H/R = 0.1$ such that the shell-averaged smoothing length per scale height is $\langle h \rangle/H \approx 0.26$ (i.e., the disk scale height is resolved by about 4 smoothing lengths). The initial disk Toomre $Q$ has a profile of $\sim r^{-3/4}$ such that the minimum $Q_{min} \approx 2.2$, at $r = r_{out}$ and $Q_{max} \approx 12.4$, at $r = r_{in}$. The simulation parameters are summarized in Table 1. The parameters are very similar to those in Fu et al. (2015b), except for the initial disk tilt. The dimensions of our parameters are all in length units of the binary separation, $a_b$, mass units of the binary mass, $M = M_c + M_p$, and time units of the orbital period of the binary, $\tau_b = 2\pi (a_{b}^{3}/G(M_c + M_p))^{1/2}$. For instance, if $a_b = 100$ au and $M = 1 M_\odot$, then both stars have masses of 0.5 $M_\odot$ and the disk extends from 2.5 to 25 au with a mass of 0.035 $M_\odot$. Our simulation lasts for about 10 binary orbits, which is about $10^4$ years. For a wider binary, the timescale for the simulation would be longer.

Our simulation tool is the PHANTOM code (Lodato & Price 2010; Price & Federrath 2010; Price 2012; Nixon et al. 2013). We use a cubic spline kernel as the smoothing kernel. The number of neighbors is roughly constant at $N_{\text{neigh}} \approx 50$. The viscosity follows the standard SPH prescription described in Monaghan (1992). A viscosity switch is implemented to reduce the artificial viscosity away from shocks (Balsara 1995; Morris & Monaghan 1997; Price & Monaghan 2004). As is standard in SPH codes, we include a nonlinear term with a coefficient $\beta_{AV} = 2$ (AV stands for artificial viscosity) in order to suppress interparticle penetration. The algorithm for the SPH implementation of self-gravity in PHANTOM is described in Price & Monaghan (2007), where the gravitational softening length is the same as the SPH kernel smoothing length (Bate & Burkert 1997). The gravitational softening length is adaptive and is approximately equal to the SPH variable pressure smoothing length. To model disk fragmentation and clump formation, we use the sink particle creation feature of the code. This converts the gas particles near a local density maximum into a new sink particle. The new sink particle does not include any gravitational softening. We follow Bate et al. (1995) in choosing the conditions for new sink particle creation. These include checking whether the total mass within the kernel radius of the particle is at least one Jeans mass, whether the velocity divergence at the particle location is negative, and whether the material that will form the new sink has both thermal energy and rotational energy less than half of its gravitational energy. To avoid checking these criteria for every local density maximum, which would greatly slow down the computation, we set a critical density above which we implement these checks. This critical density is chosen to be $500 M_{\odot} a_b^{-3}$ (note that the average disk density initially is about

| Binary and Disk Parameters | Symbol | Value |
|---------------------------|--------|-------|
| Mass of each binary component | $M_{i}/M = M_{i}/M_p$ | 0.5 |
| Binary orbital eccentricity | $e_b$ | 0 |
| Initial number of particles | $N$ | 10$^6$ |
| Initial disk mass | $M_{disk}/M$ | 0.035 |
| Initial disk outer radius | $r_{out}/a_b$ | 0.25 |
| Initial disk inner radius | $r_{in}/a_b$ | 0.025 |
| Mass accretion radius | $r_{acc}/a_b$ | 0.025 |
| Disk viscosity parameter | | 0.01 |
| Disk aspect ratio | $H/r (r = r_{in})$ | 0.1 |
| Initial disk surface density, $\Sigma \propto r^{-\gamma}$ | $\gamma$ | 1.5 |
| Initial disk inclination | $i_0$ | 70° |

Table 1

SPH Simulation Parameters for an Equal Mass Binary Star System with Total Mass $M$ and Separation $a_b$
38 $M_{\odot}^{-1}$) and an accretion radius of 0.001 $a_b$ is imposed around a newly formed sink particle. If $a_b = 100$ au and $M = 1 M_{\odot}$, then the critical density is $3 \times 10^{-10} \, \text{g cm}^{-3}$ and the accretion radius is 0.1 au.

Figure 1 shows the evolution of the disk column density up to a time of about 6.3 binary orbits. The disk orientation changes due to both orbital precession and the KL oscillations but we always show the face-on view of the disk. Because of the spatial scale of the figure, the perturber star is out of view, while the primary component is at the center of each panel. The top left panel shows that the disk is initially circular. After about three binary orbits, the disk remains circular but the perturber star has driven spiral waves in the disk. In the top right panel, which is at $t \approx 4.6 \, P_b$, the disk starts to show some eccentricity and a relatively strong one-arm spiral shock. At $t \approx 6 \, P_b$ (bottom left panel), the disk has clearly become more eccentric and the spiral shock has evolved into an arc-like shocked region with material being even more concentrated along that arc. Shortly after, at a time of $t \approx 6.28 \, P_b$ (bottom middle panel), the first clump forms (denoted by the black dot). A second clump forms shortly after this, as shown in the last panel at time $t \approx 6.3 \, P_b$. Both clumps have a mass around $3 \times 10^{-5} \, M_{\odot}$. This makes them about 10 $M_{\oplus}$ objects if we take the binary mass to be $M = 1 M_{\odot}$. They form at a radius of $r \approx 0.2a_b$, where the local disk $Q$ value is now about 1.5. By comparison, the local disk $Q$ value on the other side of the disk (also $\approx 0.2a_b$) is almost 10 times higher ($Q \approx 15$).

In Figure 2 we show the time evolution of the eccentricity and inclination of the disk at three different radii from the central star. The left column shows the run presented in Figure 1 where disk self-gravity is included. The right column, in contrast, shows a simulation without disk self-gravity (and thus no disk fragmentation) while all the other simulation parameters are the same. The left column demonstrates that disk fragmentation occurs when the disk eccentricity has grown up to almost 0.4. At this stage, the disk is still fairly early in its KL cycle, so the disk inclination has only declined slightly. If there is no fragmentation (see right column), then we see an interrupted disk KL oscillation during which disk eccentricity can grow up to about 0.7 while the disk inclination falls to about 30°. Note that in both simulations the disk remains quite flat throughout the simulation (see the similarity of disk inclination at the three different radii plotted). The KL cycle in the run that includes self-gravity starts at an earlier time than the one without self-gravity. This is because the initial driving of disk KL oscillation is very sensitive to initial model setup. Even slightly different initial conditions could lead to significantly different lengths of the pre-KL stage (as discussed in Fu et al. 2015b). Also note that the disk eccentricity in the beginning is slightly higher than zero even though we set the disk up to be circular. This is due to the fact that we assume the SPH particles are on Keplerian orbits when computing their orbital eccentricities, whereas they are actually on slightly sub-Keplerian orbits because of the disk pressure (Fu et al. 2015a, 2015b).

In Fu et al. (2015b) we presented a simulation that is the same as the one in Figure 1, except that the disk initial tilt was 50°, as opposed to 70° here. A disk with an initial tilt of 50° still undergoes KL oscillations, but the oscillation amplitudes are smaller (the peak disk eccentricity is about 0.3). There are
two reasons for the weakened KL cycle in this case. First, the disk self-gravity acts to suppress the disk KL mechanism (Fu et al. 2015b). Second, even without self-gravity, the disk KL oscillation amplitude is in general smaller for a lower initial disk tilt (see for example Figure 10 of Fu et al. 2015a). In the 50° inclination case, the disk still becomes eccentric and looks similar to the first four panels of Figure 1. However, the eccentricity growth is not strong enough for the shock front to be sufficiently dense to trigger disk fragmentation. We have also investigated two other cases that are not shown. We find that a coplanar disk does not show disk fragmentation because the KL mechanism does not operate. Furthermore, for a lower disk mass, \( M_d = 0.02 M_\odot \), the disk Q value is too high for fragmentation to occur. All of these results confirm that the disk fragmentation we see in Figure 1 is caused by disk GI that is triggered by strong eccentricity growth from a disk KL cycle. It does not happen if the disk tilt is low, which means either weak or completely no KL oscillation, at least for the disk mass (7% of the mass of the hosting star) we focus on in this study.

3. Discussion and Conclusions

In this Letter, we have shown that global KL oscillations in an inclined massive disk around one component of a binary system can facilitate the disk GI and fragmentation. For example, a disk of mass 0.035 \( M_\odot \) around a 0.5 \( M_\odot \) star can form multiple dense clumps of mass 10 \( M_\odot \) on a timescale of about 6 binary orbits. This kind of disk fragmentation can occur even when the disk does not typically have a very small Q value (e.g., \( Q > 2 \)). Thus, the same disk in isolation is fairly stable. It does require, however, that the disk initially has a high enough tilt that a strong KL-driven disk eccentricity growth can be achieved.

Three previous papers have simulated disk instability and fragmentation in binary star systems (Nelson 2000; Mayer et al. 2005; Boss 2006). They reached different conclusions regarding whether or not a binary companion helps disk fragmentation in the coplanar disk case. The discrepancy of their results is probably due to differences in the simulation techniques (SPH code or grid-based code), numerical resolution, equation of state, or binary orbit parameters (Mayer et al. 2010). Compared to these previous works, our disk is less massive (the minimum \( Q \approx 2.2 \)) and quite stable against GI without the binary companion. The disks in these three papers are just marginally stable (the minimum \( Q \approx 1.3–1.9 \)). In these papers, it generally takes a few hundred years to see disk fragmentation, if there is indeed binary companion induced disk instability. In our simulation, the disk fragments after it has gained enough eccentricity. In these papers, the disk is still quite circular at the time of fragmentation.

We have taken the disk to behave isothermally. This condition makes fragmentation occur more easily than it would in a disk that responds adiabatically due to increased resistance to compression (Gammie 2001; Rafikov 2005). Self-gravitating disks in binaries can be stabilized against fragmentation as a result of the extra heating associated with the perturbation (e.g., Mayer et al. 2010). Therefore, possible non-isothermal effects are important to include in the analysis.

We have adopted a binary separation of 100 au. A wider binary of a few hundred astronomical units or more would permit the disk to be truncated at larger radii (Artymowicz & Lubow 1994), greater than about 40 au, where a disk is more amenable to fragmentation because the optical depth is closer to unity (e.g., Hayfield et al. 2011). The KL oscillation timescale would be longer in a wider binary, but could still be shorter than the disk lifetime.

Because our code simply treats newly formed clumps as sink particles, we are not able to study the internal evolution of these dense objects and interactions between sink particles (e.g., merging of two clumps). In our study, we intentionally stop the
simulation shortly after the formation of the first one or two clumps, even though the code can still proceed with the formation of even more clumps. It is not our goal to follow the full path of the formation, evolution of any disk clump. That requires a much more sophisticated numerical investigation, which is beyond the scope of this study. Whether these disk fragments can remain as gravitationally bound objects and what their final distributions are deserve to be the subject of future investigations.

We have concentrated on KL fragmentation of protostellar disks. Other disk environments that are subject to fragmentation around a single object may experience enhanced fragmentation conditions in a non-coplanar binary environment through the KL effect. For example, fragmentation may occur in a disk around a massive black hole (Goodman 2003; Rafikov 2009). KL-induced shocks may expedite such fragmentation involving binary black hole systems with small separations where the binary mass dominates.

W.F. and S.H.L. acknowledge support from NASA grant NNX11AK61G. Computing resources supporting this work were provided by the institutional computing program at Los Alamos National Laboratory. We thank Daniel Price for providing the PHANTOM code for SPH simulations and the SPLASH code (Price 2007) for data analysis and rendering of figures.

References

Artymowicz, P., & Lubow, S. H. 1994, ApJ, 421, 651
Backus, I., & Quinn, T. 2016, MNRAS, 463, 2480
Balsara, D. S. 1995, JCoPh, 121, 357
Bate, M. R., Bonnell, I. A., & Price, N. M. 1995, MNRAS, 277, 362
Bate, M. R., & Burkert, A. 1997, MNRAS, 288, 1060
Boss, A. P. 1997, Sci, 276, 1836
Boss, A. P. 2006, ApJ, 641, 1148
Boss, A. P. 2009, ApJ, 694, 107
Evans, M. J., Ilee, J. D., Boley, A. C., et al. 2015, MNRAS, 453, 1147
Fu, W., Lubow, S. H., & Martin, R. G. 2015a, ApJ, 807, 75
Fu, W., Lubow, S. H., & Martin, R. G. 2015b, ApJ, 813, 105

Galvagni, M., Hayfield, T., Boley, A., et al. 2012, MNRAS, 427, 1725
Gammie, C. F. 2001, ApJ, 553, 174
Goodman, J. 2003, MNRAS, 339, 937
Hayfield, T., Mayor, L., Wadsley, J., & Boley, A. C. 2011, MNRAS, 417, 1839
Helled, R., Bodenheimer, P., Podolak, M., et al. 2014, in Protostars and Planets VI, ed. H. Beuther et al. (Tucson, AZ: Univ. Arizona Press), 643
Horch, E. P., Howell, S. B., Everett, M. E., & Ciardi, D. R. 2014, ApJ, 795, 60
Jensen, E. L. N., & Akesson, R. 2014, Natur, 511, 567
Kozai, Y. 1962, AJ, 67, 591
Kratter, K. M., & Lodato, G. 2016, ARA&A, 54, 271
Lidov, M. L. 1962, P&SS, 9, 719
Lodato, G., & Clarke, C. J. 2011, MNRAS, 413, 2735
Lodato, G., & Price, D. J. 2010, MNRAS, 405, 1212
Lodato, G., & Pringle, J. E. 2007, MNRAS, 381, 1287
Martin, R. G., Nixon, C., Lubow, S. H., et al. 2014, ApJL, 792, L33
Mayer, L., Boss, A., & Nelson, A. F. 2010, in Planets in Binaries, ed. N. Haghighipour (New York: Springer), 195
Mayer, L., Quinn, T., Wadsley, J., & Stadel, J. 2004, ApJ, 609, 1045
Mayer, L., Wadsley, J., Quinn, T., & Stadel, J. 2005, MNRAS, 363, 641
Meru, F. 2015, MNRAS, 454, 2529
Monaghan, J. J. 1992, ARA&A, 30, 543
Morris, J. P., & Monaghan, J. J. 1997, JCoPh, 136, 41
Nayakshin, S., & Cha, S. 2013, MNRAS, 435, 2099
Nelson, A. F. 2000, ApJL, 537, L65
Nelson, A. F., Benz, W., Adams, F. C., & Arnett, D. 1998, ApJ, 502, 342
Nixon, C., King, A., & Price, D. 2013, MNRAS, 434, 1946
Paardekooper, S., Berutau, C., & Meru, F. 2011, MNRAS, 416, L65
Papaloizou, J. C., & Savonije, G. J. 1991, MNRAS, 248, 353
Price, D. J. 2007, PASA, 24, 159
Price, D. J. 2012, JCoPh, 213, 759
Price, D. J., & Federrath, C. 2010, MNRAS, 406, 1659
Price, D. J., & Monaghan, J. J. 2004, MNRAS, 348, 123
Price, D. J., & Monaghan, J. J. 2007, MNRAS, 374, 1374
Rafikov, R. R. 2005, ApJL, 621, L69
Rafikov, R. R. 2009, ApJ, 704, 281
Rice, W. K. M., Lodato, G., & Armitage, P. J. 2005, MNRAS, 364, L56
Rice, W. K. M., Lopez, E., Forgan, D., & Biller, B. 2015, MNRAS, 454, 1940
Rogers, P. D., & Wadsley, J. 2012, MNRAS, 423, 1896
Shakura, N. I., & Sunyaev, R. A. 1973, A&A, 24, 337
Stamatellos, D. 2015, ApJ, 810, L11
Toomre, A. 1964, ApJ, 139, 1217
Tsukamoto, Y., Takahashi, S. Z., Machida, M. N., & Inutsuka, S. 2015, MNRAS, 446, 1175
Vorobyov, E. I., Zakhozhay, O. V., & Dunham, M. M. 2013, MNRAS, 433, 3256
Williams, J. P., Mann, R. K., Di Francesco, J., et al. 2014, ApJ, 796, 120