PULSED ACCRETION ONTO ECCENTRIC AND CIRCULAR BINARY

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ABSTRACT

We present numerical simulations of circumbinary accretion onto eccentric and circular binaries using the moving-mesh code AREPO. This is the first set of simulations to tackle the problem of binary accretion using a finite-volume scheme on a freely moving mesh, which allows for accurate measurements of accretion onto individual stars for arbitrary binary eccentricity. While accretion onto a circular binary shows bursts with period of $\sim 5$ times the binary period $P_{\text{bin}}$, accretion onto an eccentric binary is predominantly modulated at the period $\sim 1P_{\text{bin}}$. For an equal-mass circular binary, the accretion rates onto individual stars are quite similar to each other, following the same variable pattern in time. By contrast, for eccentric binaries, one of the stars can accrete at a rate 10–20 times larger than its companion. This “symmetry breaking” between the stars, however, alternates over timescales of order $200P_{\text{bin}}$, and can be attributed to a slowly precessing, eccentric circumbinary disk. Over longer timescales, the net accretion rates onto individual stars are the same, reaching a quasi-steady state with the circumbinary disk. These results have important implications for the accretion behavior of binary T Tauri stars and supermassive binary black holes.

Key words: accretion, accretion disks – binaries: general – black hole physics – stars: pre-main sequence

1. INTRODUCTION

Spectroscopic T Tauri star binaries can exhibit quasi-periodic photometric oscillations known as “pulsed accretion” (Jensen et al. 2007; Muzerolle et al. 2013; Bary & Petersen 2014). This variability is believed to arise from the complex accretion streams delivered onto the young stars from a tidally truncated circumbinary disk (e.g., Artymowicz & Lubow 1996). Similar circumbinary disks may also exist around supermassive binary black holes (SMBBHs), but the short periods of binary T Tauri stars (BTTSs) offer an unparalleled coverage of the time domain as the system can evolve over several orbits during observations and even up to hundreds of orbits between different observing campaigns. This makes BTTSs ideal laboratories for circumbinary accretion physics, with direct implications for binary star and circumbinary planet formation, and with applications that extend to SMBBHs.

The complexity of circumbinary accretion flow requires direct hydrodynamical simulations. Several computational approaches have been taken to address this problem (Artymowicz & Lubow 1996; Günther & Kley 2002; MacFadyen & Milosavljevic 2008; Cuadra et al. 2009; Hanawa et al. 2010; de Val-Borro et al. 2011; Roedig et al. 2012; Shi et al. 2012; D’Orazio et al. 2013; Pelupessy & Portegies Zwart 2013; Farris et al. 2014; Lines et al. 2015), ranging from Lagrangian methods to Eulerian ones on polar and Cartesian grids. Of these, only a subset has been able to simulate the accretion flow onto the individual stars, or to include eccentricity in the binaries. To date, only four studies have attempted both; two of these using grid-based approaches (Günther & Kley 2002; de Val-Borro et al. 2011), and two using smoothed particle hydrodynamics (SPH; Artymowicz & Lubow 1996; Dunhill et al. 2015).

Considering eccentric binaries in simulations is essential for pulsed accretion, as accretion luminosity is likely to depend on the orbital phase (Basti et al. 1997; Huerta et al. 2005; Jensen et al. 2007; Bary et al. 2008). Although high eccentricities in BTTSs are common—the binaries AK Sco, DQ Tau, and UZ Tau E have eccentricities of $e_b = 0.47$, 0.56, and 0.29, respectively (Andersen et al. 1989; Mathieu et al. 1997; Prato et al. 2002)—accurate simulation of circumbinary accretion onto eccentric pairs remains a challenge. In this work, we present the first simulation results of circumbinary accretion using the moving-mesh code AREPO (Springel 2010). Unlike other implementations of finite-volume or finite-difference schemes for computational gas dynamics, the accuracy of AREPO does not depend on the value of $e_b$, as its space-discretization strategy is carried out via an unstructured mesh that moves with the local velocity of the flow. Thus, being a quasi-Lagrangian method, AREPO can naturally concentrate the resolution around the individual stars, resolving circumbinary disks with sufficient accuracy. At the same time, being fundamentally a “grid code,” AREPO does not suffer from some of the limitations inherent to other Lagrangian-like schemes such as SPH, including poor shock-capturing performance, suppression of instabilities, and spurious surface tension (for a suite of comparisons of AREPO and SPH, see Sijacki et al. 2012).

2. NUMERICAL METHODS

2.1. Moving-mesh Hydrodynamics

We run two-dimensional, non-self-gravitating hydrodynamic simulations of viscous circumbinary accretion disks (CBDs) using AREPO with a time-explicit integration scheme for the Navier–Stokes terms (Muñoz et al. 2013). The computational domain is divided into Voronoi cells, distributed in a quasi-polar fashion with logarithmic spacing in radius, following the accretion disk setup of Muñoz et al. (2014). Cells initially cover the radial range from $R = a_0(1 + e_b)$ to $R = R_{\text{out}} = 70a_0$ (where $a_0$ is the binary semimajor axis) but are allowed to viscously evolve toward $R = 0$. Resolution elements in the inner CBD ($R < 10a_0$) are distributed into 180 radial zones and 600 azimuthal zones, while resolution in the outer CBDs ($10a_0 \leq R \leq 70a_0$) consists of 120 radial zones and 300 azimuthal zones. At $R_{\text{out}}$, we impose inflow boundary
conditions of steady accretion $M_b$, assuming that at these distances the disk is axisymmetric and the central potential Keplerian. The binary is represented by a prescribed rotating potential:

$$\Phi(r) = -\frac{GM_b}{|r - r_1|} \left[ \frac{(1 + q_b)^{-1}}{|r - r_1|} + \frac{q_b(1 + q_b)^{-1}}{|r - r_2|} \right], \quad (1)$$

where $q_b = M_2/M_1$ is the binary mass ratio and $M_b = M_1 + M_2$ is the total mass. The individual stellar positions are $r_1(t) = q(1 + q)^{-1}r(t)$ and $r_2(t) = -(1 + q)^{-1}r(t)$, where the relative position vector $r(t) = a_0 \cos \epsilon - e_0 \sin \epsilon$. The eccentric anomaly $E(t)$ is obtained by solving Kepler’s equation (e.g., Danby 1988). The potential around each star is softened, the softening length is set to $s = 0.025 a_0$.

The binary components are allowed to “accrete” (although their dynamical masses are held constant). Gas is drained from cells located within a distance of $a_{acc} = 0.8 a_0$ from each star. The draining is carried out as a simple “open-boundary” condition, meaning that cells that are located within the accretion region are instantaneously drained (Muñoz et al. 2015). The accreted mass $M_i$ is stored at every major time step, and accretion rates $M_i$ and $M_2$ are computed by central-difference differentiation of $M_i(t)$. Note that, for BTTSs with semimajor axis $a_0 \sim 0.2$ au (Jensen et al. 2007), we have $a_{acc} \approx 0.004 a_1 R_\odot$, sufficient to resolve the accretion down to the stellar surface. On the other hand, for an SMBBH, the “true” accretion radius (e.g., the innermost stable circular orbit) is $< a_{acc}$, and the value of $M_i$ should be interpreted with caution (see Section 3.4).

As the outer CBD evolves, resolution is maintained roughly constant via de-refinement and refinement operations (Springel 2010). Within $R = a_0 (1 + e_0)$, the resolution criterion is switched over from “volume-based” to “mass-based,” in which there is a “target mass” $m_{gas}$ enforced for all cells (Springel 2010). The transition between mass-based and volume-based resolution is kept smooth by controlling the volume difference between contiguous cells (Pakmor et al. 2013). For our lowest-resolution runs, $m_{gas} = 6.3 \times 10^{-7} \Sigma_0 a_0^2$, where $\Sigma_0$ is the scaling of the initial disk surface density profile (see Section 2.2 below). In this region, there is a minimum permitted volume, $\pi s^2/20$, where $s$ is the softening parameter.

The equation of state is “locally isothermal,” $P = \Sigma c_s^2(r)$, where the sound speed is a function of position only (e.g., Farris et al. 2014): $c_s^2(r) = -h_0^2 \Phi(r)$, where the aspect ratio $h_0$ is a global constant. When $|r_1| \gg |r|$, $|r_2|$, then $c_s^2 \approx h_0^2 G M_b/|r|$; and when $|r - r_1| \ll |r - r_2|$, then $c_s^2 \approx h_0^2 G M_b/|r - r_1|$.

Finally, the kinematic viscosity $\nu$ follows an $\alpha$-viscosity prescription (Shakura & Sunyaev 1973), in which $\nu = \alpha c_s^2/\Omega(r)$, where $\Omega(r)$ is a function that reduces to $(GM_b/R)^{1/2} \nu$ far from the binary and to $(GM_b/(r - r_1))^{1/2} \nu$ close to each star.

Throughout this work, we fix the parameters $q_b = 1$ and $h_0 = \alpha = 0.1$, while varying the binary eccentricity $e_b$.

2.2. Initial Setup

Knowing that the outer disk is in steady-state accretion, we “guess” an initial surface density profile $\Sigma(R)$ that includes a central cavity but that, at large radii, behaves as $\Sigma \propto M_b/\nu \propto R^{-1/2}$. Thus, we adopt the initial CBD surface density profile:

$$\Sigma(R, t = 0) = \Sigma_0 \left( \frac{R}{R_{cav,0}} \right)^{-p} \exp \left[ -\left( \frac{R}{R_{cav,0}} \right)^{-\xi} \right], \quad (2)$$

where $p = 1/2$ and $R_{cav,0}$ and $\xi$ characterize the extent and steepness of the tidal cavity around the binary. In this work, we choose $R_{cav,0} = 5 a_0$ and $\xi = 4$. Steady state at the onset of the simulation is guaranteed for $R \gg R_{cav,0}$ by construction, but Equation (2) is still an imperfect initial condition at intermediate radii, and a long integration time may be needed to relax the initial conditions for all $R$.

The initial condition is completed by specifying a rotation curve

$$\Omega^2 = \frac{GM_b}{R^3} \left[ 1 + \frac{3}{4} \frac{a_0}{R} \left( \frac{1}{1 + q_b} \right)^2 \left( 1 + \frac{3}{2} e_b^2 \right) \right] + \frac{1}{R \Sigma} \frac{d P}{d R},$$

which includes the quadrupole component of the potential and the contribution of the pressure gradient, and by specifying a radial velocity profile $v_R(R)$. Assuming a standard thin accretion disk, we impose

$$v_R = \frac{1}{R \Sigma} \frac{\partial}{\partial R} \left( \nu \Sigma R^3 \frac{d \Omega}{d R} \right) \left( \frac{d}{d R} (R^2 \Omega) \right)^{-1},$$

which in turn specifies the accretion rate profile

$$\dot{M}(R) = -2 \pi R v_R(R) \Sigma(R).$$

Note that this initial $\dot{M}(R)$ starts converging toward $M_b$ only beyond $R \geq 20 a_0$ (at $t = 0$, $\dot{M}$ equals $1.41 M_b$ and $1.08 M_b$ at $R = 10 a_0$ and $15 a_0$, respectively), and thus the disk is not started with a strictly steady accretion profile. Unless stated otherwise, we initially evolve the system for $200 P_b$ (where $P_b$ is the binary orbital period) and study the subsequent evolution for an additional 600 binary orbits. This initial integration time corresponds to two viscous times at $6 a_0$ or $10 a_0$, where the viscous time $t_v$ is defined for $\nu \propto R^{1/2}$ (Lynden-Bell & Pringle 1974) as

$$t_v = \frac{4 R^2}{9 \nu} = \frac{2 P_b}{9 \pi a_0^2} \left( \frac{R}{a_0} \right)^{3/2}.$$

For 1950 orbits, we evolve the disk using an open (diode-like) boundary on a set of controlled cells placed on a ring at $R_{in} = a_0 (1 + e_0)$. At $t = 1950 P_b$, boundary cells are “released,” allowing them to fill in the cavity and form accretion disks around the individual stars. At $t \geq 2000 P_b$, we expect the CBDs within...
5\theta_0 to be (on average) fully relaxed. We aim to reach a “relaxed state” inside the cavity as well, for which \( \langle M_1 + M_2 \rangle \approx \langle M(R) \rangle \approx \) constant for a wide range in \( R \), after averaging over some time interval \( T \). The outer disk \( R > 4a_0 \) is in steady state by construction. However, there is an intermediate region, with \( t_v > 2000P_b \), that has had no time yet to relax (see below).

2.3. Long-term Disk Relaxation

Inspection of the initial condition reveals that \( M(R, t = 0) \) coincides with \( M_0 \) to within 1% only for \( R > 42a_0 \). This is just an artifact of the initial condition. To guarantee \( M(R) \approx M_0 \) across all radii, we would need to evolve the system for \( \sim 20,000 P_b \) (or \( t_v \approx 40a_0 \), Equation (6)), a daunting task for the simulation work presented here. Instead, after some integration time \( t_{\text{int}} \), the system has only reached relaxation within a “relaxation radius” \( R_{\text{rel}} = a_0 \left( \frac{9}{2} \pi \alpha h_0 (t_{\text{int}}/P_b) \right)^{2/3} \) (from setting \( t_{\text{int}} = t_v(R_{\text{rel}}) \) in Equation (6); see Rafikov 2016). After \( t = 2000P_b \), \( R_{\text{rel}} \approx 3.3a_0 \), which is \(<42a_0 \), but sufficiently large for the disk to be nearly axisymmetric outside that radius. Without reaching global relaxation, we have found that the disk within \( R_{\text{rel}} \) receives a gas supply from the partially relaxed portion of the disk \( R_{\text{rel}} \approx R \approx 2R_{\text{rel}} \) at a rate \( M_{\text{out}} \) slightly larger (by about 10%) than \( M_0 \), as a result of the initial condition. This accretion “excess” cannot be removed with only 2000 orbits of integration time.

3. SIMULATION RESULTS

3.1. Accretion Flows in the Circumbinary Cavity

Figure 1 shows several snapshots of the density field for accretion onto a circular binary (top) and an eccentric one (bottom). The tidal “cavity” around the binary is asymmetric and noncircular, making it difficult to identify an unambiguous cavity “radius.” Within the cavity, flow is complex and transient, dominated by accretion streams, which show significantly more structure (and unsteadiness) in the eccentric case. In the circular case, the “streamers” rotate with the binary, while this is not the case when \( e_b \approx 0 \). Similarly, the cavity shape and contrast nearly repeat themselves every half orbit when \( e_b = 0 \), while no such symmetry is observed when \( e_b = 0.5 \). In both cases, a well-resolved, rotationally supported CSD forms around each star, showing that gas does not free-fall directly from the streamers into the accretion region. In the eccentric case, the CSDs are severely truncated at pericenter, but at apocenter, they have re-expanded thanks to the newly delivered gas from the CBD.

The symmetries of the \( e_b = 0 \) case are also evident from the properties of the disks: both CSDs are similar in size, density, and morphology (including \( m = 2 \) spiral patterns in each). By contrast, the CSDs in the \( e_b = 0.5 \) case show differences in surface density, implying that the members of the

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Footnotes:

[1] The annulus that supplies gas to the inner CBD can be roughly estimated by integrating the traveled distance from \( R_{\text{out,eff}} \) down to \( R_{\text{rel}} \) at a velocity of \( v_R = \frac{\sqrt{2}}{\rho} (\frac{1}{R} - \frac{1}{R_{\text{out,eff}}})^{1/2} \), giving \( R_{\text{out,eff}}/R \approx 1.17 \) if \( t_{\text{int}} = 2000P_b \). From Equations (2) and (5), we have \( M(R) = 14.7a_0, t = 0 \approx 1.3M_0 \).
equal-mass binary accrete at different rates. We will further address this “disk disparity” in Section 3.3.

3.2. Accretion Rates

A schematic description of the circumbinary accretion process is the following: there are three accretion disks—the CBD and the two CSDs—that evolve viscously, but are connected via fast accretion bursts owing to the (unstable) tidal streams launched at the inner edge of the CBD. If $e_b = 0$, the stars may collect new material at apocenter, when they are closest to the CBD inner edge. Subsequently, the incoming material is viscously transported inward within the CSD, eventually accreting onto the stars at some later phase in the orbit. The (presumably) much slower rate at which material is transported onto the accreting objects relative to fast deposition of material to the outer edge of the CSDs turns the CSDs into “buffers.” The buffers damp the fast oscillations in $\dot{M}$ present in the circumbinary cavity before they reach the stars.

Although compelling, this idealized depiction is clearly too simplistic in light of the simulation results of Figure 1. We compute $\dot{M}(R)$ (Equation (5)) at different radii\(^8\) of the CBD. Figure 2 shows $\dot{M}(R)$ at $R = (1 + e_b)$, 3, 5, 7, and $10 \times a_0$ as a function of time for $e_b = 0$ (left) and $e_b = 0.5$ (right). In particular, $R = a_0(1 + e_b)$ is where the innermost boundary would be located in a polar-grid simulation. The total accretion onto the central masses $M_{\text{bin}} \equiv M_1 + M_2$ is shown on top.

\(^8\) We compute $\dot{M}(R, \phi) = -2\pi R v_R(R, \phi) \Sigma(R, \phi)$ for all Voronoi cells in the vicinity of radius $R$ and then take a weighted azimuthal average.
Accretion rates are normalized to a reference value $\dot{M}_{\text{ref}}$, which is the accretion rate at a radius where the disk becomes axisymmetric. We measure the mean $\langle M \rangle$ over a period $T = 60P_b$ in each panel (see figure caption). In the $e_b = 0$ case (Figure 2, left panels), there is a clear accretion modulation with period $\sim 5P_b$ (roughly the Keplerian period at $R = 3a_b$), observed in $\dot{M}_{\text{ref}}$ (top panel) as well as in the CBD out to $R = 3a_b$. Right at the putative edge of the CBD disk ($\sim 3a_b$), the evolution of $M$ turns significantly more complex (a quasi-periodicity of $5P_b$ is still present) and extremely variable in amplitude (going from $-40$ to $+20 \dot{M}_{\text{ref}}$). For $R < 3a_b$, modulation of $M$ is dominated by a $\sim 0.5P_b$ component superposed to major bursts that repeat every $\sim 5P_b$. The periodicity at $5P_b$ could not capture this sign-changing behavior. Orazio et al. (2013; Farris et al. 2014) have attributed these bursts to an overdense “lump” that forms at the rim of the cavity and gets periodically “flung” onto the binary (see Figure 1, top panels). Note that at $R = 1a_b$—where a polar-grid code would place the outflow computational boundary—$M$ is always positive, and thus artifacts introduced by diode-like boundary conditions (not allowing for material with $v_R > 0$ to enter the domain) are minimal. The two top panels of Figure 2 show good qualitative agreement with each other, the differences being (1) a delay in the time of the accretion burst to reach the stars and (2) a reduction of the amplitude of the variability; these differences are consistent with the buffering nature of a viscous disk.

The $e_b = 0.5$ case (Figure 2, right panels) shows much more complex $M$ variability outside $R = 3a_b$. By contrast, for $R < 3a_b$, variability seems simpler than around a circular binary. In particular, $\sim 1P_b$ is the dominant modulation period, although trends of periods longer than $60P_b$ are also noticeable. The amplitude of the oscillations in $\dot{M}_{\text{ref}}$ (Figure 2, top right panel) can be as high as $5\dot{M}_{\text{ref}}$, in contrast with $\sim 1.5\dot{M}_{\text{ref}}$ for the $e_b = 0$ case. Another striking difference from the $e_b = 0$ case is the value of $M$ around the binary. The imaginary boundary at $R = a_b (1 + e_b)$ (second right panel from top) shows alternating negative and positive values of $M$. This is consistent with the appearance of shocks inside the cavity (Figure 1, bottom panels) as a result of the convergence of inflowing and outflowing streams. Evidently, a diode-like boundary placed at $R = a_b (1 + e_b)$ could not capture this sign-changing behavior. By contrast, $M$ at this location is always positive around a circular binary (Figure 2, left, second from top).

### 3.3. Individual Accretion Rates and Correlation with Periastron Separation

In Figure 3 we show accretion onto the individuals stars $M_1$ and $M_2$ over 160 binary orbits. For the $e_b = 0$ case (left panels), the symmetry between the primary and secondary is remarkable, as it is expected for $q_b = 1$. This is at odds with the results of Farris et al. (2014), which show a mild “symmetry breaking” in $M_1$. Both $M_1$ and $M_2$ show the bursty nature of the two top left panels of Figure 2, although they do not perfectly lie on top of each other; instead, one star undergoes an accretion burst before its companion. The lag between the two bursts is about a half orbit, although the sign of the lag alternates on each major burst.

For $e_b = 0.5$ (right panels of Figure 3), the symmetry breaking between $M_1$ and $M_2$ is evident. Despite having $q_b = 1$, over the
first ~90 orbits, $M_1$ (blue) is 10–20 times larger than $M_2$ (red). Interestingly, after 100 orbits, this behavior switches to $M_2 > M_1$, only to switch back to $M_1 > M_2$ at $t = 2300P_b$ (not shown). Over timescales of ~600$P_b$, we have that $M_1 \approx M_2$, recovering—in a time-averaged sense—the symmetry that is to be expected when $q_b = 1$. The reason for this dramatic difference between $M_1$ and $M_2$ must originate in a symmetry breaking in the CBD itself.

If the CBD is eccentric, the relative longitude of pericenter $\omega_d - \omega_b$ (where $\omega_d(R)$ specifies the orientation of a given elliptical portion of the CBD disk) will determine the timing of mass transfer from the CBD to the binary. In principle, an eccentric disk should precess around the binary at a rate $\dot{\omega}_d$, implying that, if one of the accreting objects is benefited by an increased $M$ at any given time, at some later time preferential accretion should be reversed. A relevant precession rate is that of the inner rim of the CBD, at a radius of $R_{cav} \sim 2 \sim 3 a_b$. In the limit of a pressure-less particle disk, the secular apsidal precession rate around an eccentric binary is

$$\dot{\omega}_d \approx \frac{3\Omega_b}{4} \frac{q_b}{(1 + q_b)^2} \left(1 + \frac{3}{2} e_b^2\right) \frac{d_b}{R}^{7/2},$$

$$\sim 0.006 \Omega_b \left(\frac{3 a_b}{R}\right)^{7/2},$$

which corresponds to a precession period of a few hundred $P_b$ at $R \sim 3 a_b$. This precession period roughly coincides with the period of alternation of dominant accretion shown in Figure 3. In principle, a nonaxisymmetric potential could also induce oscillations in the CBD eccentricity $e_d$, which in turn could contribute to the variability of the stellar accretion rates (e.g., Lubow & Artymowicz 2000). Note, however, that $e_d = 0$ to quadrupole order in the perturbing potential. To octupole order, the change in eccentricity is given by (e.g., Lubow & Artymowicz 2000; Moriwaki & Nakagawa 2004)

$$\dot{e}_d \approx -\frac{15}{16} \frac{q_b}{(1 + q_b)^3} \frac{a_b}{R} \left(\frac{a_b}{R}\right)^{9/2} \sin(\omega_d - \omega_b),$$

which is identically zero for the simulation work presented here ($q_b = 1$). Thus, any pulsating behavior in $e_d$ when $q_b = 1$ must be a consequence of resonant excitation and gas dynamical effects, and not due to secular dynamics. Future work will take a deeper look into the properties of precessing eccentric CBDs (R. Miranda et al. 2016, in preparation).

Recently, Dunhill et al. (2015) have reported a similar, long-term alternating behavior in the accretion rate onto eccentric binaries using SPH simulations. Interestingly, their simulations do not exhibit the high-frequency oscillations in accretion onto the individual stars reported here and in Farris et al. (2014). The reason for this discrepancy could lie in the lack of circumstellar disk formation in such SPH simulations, despite the high resolution and the use of a similarly sized accretion radius ($0.03 a_b$) to the one used in this work ($0.02 a_b$).

In Figure 4, we show a portion of Figure 3 ($e_b = 0.5$ case, right panels) overlaid with the binary separation [n $\sim n$]. In the case where $M_{bin}$ is dominated by $M_1$ (blue curve), accretion peaks before pericenter passage, exhibiting a minor second peak exactly at pericenter. This is in partial agreement with the simulation results of Günther & Kley (2002) and de Val-Borro et al. (2011), although the accretion burst peaks noticeably before pericenter, and the burst duration spans a significant fraction of the orbital period.

### 3.4. Tidal Torques and the Effect of Increased Resolution

We now examine the buffering nature of the CSD discussed above (Section 3.2). When the accretion radius $r_{sec}$ (Section 2.1) is much smaller than the size of the CSD $R_{cav}$, the accretion time $t_{sec,cav}$ within a CSD is roughly the viscous time at the disk edge $t_{sec,cav} = 2P_b/(9\pi\alpha_b^2) (R_{cav}/a_b)^{3/2}/1 + q_b$. This timescale enables the damping of fast modulations and sets a delay between the time of gas deposition onto the CSD and the time of actual accretion onto the stars. With an estimate of $R_{cav},^{10} we have $t_{sec,cav} \approx 23 P_b$. This time may be short enough to enable the accretion burst of pericenter ~5$P_b$ to reach the stars, but it is perhaps too long to allow for the persistence of the high-frequency oscillations (periods ~1$P_b$ and shorter; Figures 2 and 3, left panels). However, the fast oscillations in $M_{bin}$ might be explained by the enhanced mass transport via “spiral shocks” launched by an external potential (Rafikov 2016). In a circum-primary frame (primed coordinates), this external potential on the CSD is due to the secondary (e.g., Miranda & Lai 2015):

$$\Phi_{sec}(r', t) \approx \frac{GM}{a_b} \frac{q_b}{1 + q_b} \frac{1}{4} \frac{r'^2}{a_b^2} [1 + 3 \cos(\phi' - \Omega_b t)],$$

where we have dropped a constant term. The time-dependent term can excite $m = 2$ density waves in the circum-primary disk (Figure 1). These features are similar to the spiral shocks observed in simulations of disks inside large mass ratio binaries (see Dong et al. 2016; Ju et al. 2016, for recent examples). Spiral shocks can transport angular momentum and enhance the accretion rates in disks (Goodman & Rafikov 2001; Rafikov 2016), potentially making it the dominant mechanism for mass transport over local viscosity. Following Rafikov (2016), one can roughly define a shock-driven Shakura–Sunyaev parameter $\alpha_{shock} = (m/r)\psi \bar{\psi}$, where $\psi$ is a measure of irreversible heating (which is instantaneously balanced by heat extraction in isothermal disks),

$$\psi = [0.49q_b^{-2/3}/(0.6q_b^{-2/3} + \ln(1 + q_b^{-1/3})].$$

9 If the condition $t_{sec} \ll t_{cav}$ is not satisfied, a more general expression for the accretion rate is $t_{sec,cav} = t_{cav}(1 - (t_{sec}/t_{cav})^{1/2})$.

10 We replace $R_{cav}$ with the Eggleton approximation of the Roche radius (Eggleton 1983):

$$(R_{cav}/a_b) = 0.49q_b^{-2/3}[0.63q_b^{-2/3} + \ln(1 + q_b^{-1/3})].$$
taking values \( \psi_0 \sim 0.06-0.5 \) for compression levels of \( \Delta \Sigma/\Sigma_0 \sim 1-3 \). Thus, spiral shocks imply \( \alpha_{\text{shock}} \sim 0.04-0.3 \), introducing moderate to large modifications to the accretion rates that one would expect from viscous evolution alone. This can help explain why the fast oscillations in \( \dot{M} \) measured at \( R = a_b(1 + e_b) \) can survive the presumed buffering effect of the CSD.

Note that our accretion routine is somewhat resolution dependent (cells are drained depending on their location, regardless of their total mass content; see Muñoz et al. 2015). For a given \( M_0 \), the average number of cells being accreted in an interval \( \Delta t \) is \( N_{\text{acc}} = (M_0/m_{\text{gas}})\Delta t \), with a “signal-to-noise ratio”\(^{11} \) of \( \sqrt{N_{\text{acc}}} \approx 200\sqrt{\Sigma_0/\Delta t} \). For \( \Delta t \sim 0.05P_b \), the uncertainty in the measured \( M_{\text{bin}} \) is \( \sim 1\% \), small enough to be confident in the general features of Figure 3, but large enough to justify a convergence study of the high-frequency modulations (see Pakmor et al. 2016). Figure 5 (top panels) shows the circum-primary density field at three different gas resolutions (while keeping \( r_{\text{acc}} \) and the softening length \( s \) fixed), confirming the prevalence of \( m = 2 \) spiral arms. The bottom panel of Figure 5 shows \( M_{\text{bin}} \) at the same three resolutions, confirming the major accretion burst and that the rapid oscillations are real and likely a result of a time-dependent forcing (Equation (10)). Note that the strength of this forcing in Equation (10) is negligible for \( r' \ll a_b \); thus, one can expect the influence of the companion and the amplitude of the resulting spiral shocks to be drastically reduced as \( r_{\text{acc}} \) is made smaller. We repeat the resolution experiments of Figure 5 (not shown), this time decreasing \( r_{\text{acc}} \) and \( s \) in addition to \( m_{\text{gas}} \). We find that the fast modulations are progressively damped out; and for very small \( r_{\text{acc}} \), only the major accretion bump survives, as the accretion (buffering) time within the CSD is not long enough to entirely erase it.

4. SUMMARY AND IMPLICATIONS

We have presented two-dimensional, viscous flow simulations of circumbinary disk accretion for the first time using a finite-volume method on a freely moving Voronoi mesh. Previous simulations based on structured moving grids were restricted to circular binaries (e.g., Farris et al. 2014). Using AREPO, we can robustly simulate accretion onto arbitrarily eccentric binaries, without the constraints imposed by structured grids. In our simulations, we are able to follow the mass accretion through a wide radial extent of the circumbinary disk, leading to accretion onto individual members of the binary via circumstellar disks.

Our simulations have revealed dramatic differences between the accretion behavior of circular and eccentric binaries:

1. In agreement with previous studies (e.g., Shi et al. 2012; D’Orazio et al. 2013; Farris et al. 2014), we find that accretion onto equal-mass, circular binaries exhibits quasi-periodic variabilities with a dominant period of \( \sim 5P_b \) (where \( P_b \) is the binary period), corresponding to the orbital period of the innermost region of the circumbinary disk. By contrast, accretion onto eccentric binaries displays larger-amplitude variabilities dominated by pulses with periods of \( \sim 1P_b \) (see Figure 2).

2. For equal-mass circular binaries, we find that the accretion rates onto individual stars are quite similar to each other, following an essentially identical accretion pattern in time (Figure 3, left panels). This result differs

\(^{11}\) The imposed accretion rate \( M_0 \) and mass resolution are related by \( M_0 = 3\sqrt{5}\pi c_0 h_0^2 (R_{\text{acc}}/5a_b)^{1/2} (m_{\text{gas}}/0.03)(6.7 \times 10^{-3}) \).
from the simulations by Farris et al. (2014), which produced an appreciable disparity between the individual stellar accretion rates. By contrast, we find that accretion onto eccentric binaries exhibits strong symmetry breaking: for a period of time lasting ~200 $P_b$ (which corresponds to the apsidal precession period of the innermost region of the circumbinary disk), one of the stars can accrete 10–20 times more than the companion (Figure 3, right panels). This disparity alternates over timescales of ~200$P_b$, such that the long-term accreted masses onto individual stars are the same.

In addition to using a novel moving-mesh code (AREPO) that resolves the binary-disk system over a large dynamical range, an important feature of our study is that we carry out our simulations for a sufficiently long time (thousands of binary orbits) and with a proper initial setup. The inner circumbinary disk and the individual circumstellar disks reach a quasi-steady state in which the time-integrated mass accretion is the same across different regions of the system. The ability to reach quasi-steady state gives us confidence that the pulsed accretion behavior uncovered in this paper is not the result of artificial initial conditions.

Our results can be compared to the observations of pulsed accretion in BTTSs (Jensen et al. 2007; Muzelle et al. 2013; Bary & Petersen 2014) and can shed light on the origin of the quasi-periodic variability (in broadband photometry and near-IR line fluxes) observed in these systems. In our simulations, accretion onto eccentric binaries peaks before and during pericenter but never at apocenter. This appears to contradict the observation of BTTS DQ Tau (with $e_b = 0.56$), which exhibits flaring events during apocenter (Bary & Petersen 2014). This apparent discrepancy can be easily understood once we recognize that the size of the accretion region $r_{acc}$ can strongly affect the measured variability of accretion rates (Section 3.4). In the case of DQ Tau, pericenter passage of the binary (at the separation of $r_p = a_b(1 - e_b) \approx 0.05$ au; Mathieu et al. 1997) would limit the size of circumstellar disks to be less than $\sim r_p/3 \approx 0.02$ au, i.e., about 2–3 pre-main-sequence stellar radii, making circumstellar disk accretion irrelevant (especially if the stellar magnetospheres are indeed colliding at periastron; Salter et al. 2010), with accretion proceeding almost directly from the streamers to the stars. In addition, it is possible that the shocked gas responsible for line emission is not strictly confined to the stellar photospheres (Calvet & Gullbring 1998), but located elsewhere in the circumbinary cavity (e.g., Bary & Petersen 2014). Indeed, our eccentric binary simulations do show that shocks appear as material is swung out from the edges of the circumstellar disk at each close passage; this outflowing material meets the inflowing accretion streams from the circumbinary disk. As the disk flow is highly supersonic, the eccentric accretion streams could shock against material at a relative Mach number of $\mathcal{M} \sim 10$. We plan to explore the observational signatures of these shocks in future work.

Finally, although we have focused on accretion onto pre-main-sequence binaries in this paper, our simulations also have implications for accretion onto SMBBHs. In Section 3.4 we have discussed how the size of the accreting region $r_{acc}$ can affect the variability of accretion rates, such that when $r_{acc} \to 0$, modulations of accretion on timescales much shorter than the circumstellar disk viscous time are damped out. Since we expect $r_{acc} \ll a_b$ for SMBBHs, the individual black holes in a binary would accrete at the nominal supply rate, suppressing fast variability. If this is the case, the mechanism behind the observational hints of photometric variability of SMBBHs would most likely be due to Doppler beaming, as suggested by O’Rorio et al. (2015), rather than to gas dynamics within the circumbinary cavity.

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