1. INTRODUCTION

Mass transfer from an evolved donor star to a companion star is a common outcome of stellar evolution in binary systems. In close binaries with periods of 1–2 yr or less, the donor can fill its tidal surface and transfer mass rapidly into a massive disk close binaries with periods of 1–2 yr or less, the donor can fill its tidal lobe. In both cases, long-lived disks are possible when the donor star does not fill its tidal lobe.

Theoretical studies of the physical properties of wind-accretion disks provide an important framework for interpreting observations of these systems. The replenishment of newly accreted material and the formation of a disk may also have important implications for the formation and evolution of planetary systems. A newly formed disk may interact with pre-existing planets, leading to the re-growth and/or migration of these planets, or possibly the formation of new planets and/or debris disks (Perets 2010, 2011).

Several previous studies explored the formation of wind-accretion disks through smoothed particle hydrodynamics and AMR simulations (Mastrodemos & Morris 1998; de Val-Borro et al. 2009). Here, we use a different simplified analytical/numerical approach typically used for the study of disks in shorter period binaries and in protostellar/protoplanetary systems. This approach allows us to follow the long-term evolution of the disks, which is more difficult in hydro simulations, and to compare the derived physical structure with protoplanetary disks. We focus on the accretion from low-mass companions in wide binaries (3–100 AU), which evolve through wind accretion.

We begin with a description of the basic conditions in wind-accreting binaries, followed by a detailed explanation of the methods we use to follow their evolution. We then describe results of the wind-accretion disk structure followed by a discussion and summary.

2. WIND ACCRETION IN WIDE BINARIES

To understand the conditions required for the formation of circumstellar disks from wind-accreted material, we follow Soker & Rappaport (2000). For a disk to form, the specific angular momentum $j_a$ of the accreted material, must exceed $j_z = (GM_*/R_*)^{1/2}$, the specific angular momentum of a particle in a Keplerian orbit at the equator of the accreting star of radius $R_2$. For accretion from a wind, the net specific angular momentum of the material entering the Bondi–Hoyle accretion radius, $R_\text{BH}$, is $j_{BH} = 0.5(2\pi/P)R_\text{BH}^2$ (Wang 1981), where $P$ is the orbital period. The actual accreted specific angular momentum for high Mach number flows is $j_a = \eta j_{BH}$, where $\eta \sim 0.1$ and $\eta \sim 0.3$ for isothermal and adiabatic flows, respectively (Livio et al. 1986). For a binary composed of a main-sequence
(MS) accretor with mass $M_2$ and a mass-losing asymptotic giant branch (AGB) star with mass $M_1$, requiring $j_u > j_f$ leads to a simple condition for disk formation

$$1 < \frac{j_u}{j_f} \simeq 1.2 \left( \frac{\eta}{0.2} \right) \left( \frac{M_1 + M_2}{2.5 M_\odot} \right) \left( \frac{M_2}{M_\odot} \right)^{3/2} \left( \frac{R_2}{R_\odot} \right)^{-1/2} \left( \frac{a}{100 \text{ AU}} \right)^{-3/2} \left( \frac{v_r}{10 \text{ km} \text{s}^{-1}} \right)^{-4}, \quad (1)$$

where $R_2$ is the radius of the accreting star, $a$ is the semimajor axis of the binary, and $v_r$ is the relative velocity of the wind and the accretor.

Adopting $R_2 \sim R_\odot$, disks can form in binaries with $a \lesssim 10$–100 AU for $M_1 \approx 1$–10 $M_\odot$ and $v_r \approx 5$–20 km s$^{-1}$. This range in $a$ spans the peak of the distribution of binary star separations (Duquennoy & Mayor 1991). Thus, many binaries contain wind-accretion disks during the late stages of their evolution.

Red giant branch (RGB) and AGB stars lose mass in low-velocity winds at typical rates of $\dot{M}_w \sim 10^{-8}$–$10^{-5}$ $M_\odot$ yr$^{-1}$ (e.g., Taylor & Seaquist 1984; Kenyon 1986; Olofsson et al. 2002; Ramstedt et al. 2006). To estimate the fraction of this material accreted by the MS companion, we adopt the Bondi–Hoyle rate for an isotropic wind from the mass-losing star,

$$\dot{M}_a = \left( \frac{R_a}{2a} \right)^2 \dot{M}_w,$$ \quad (2)

where $R_a$ is the Bondi–Hoyle accretion radius

$$R_a = \frac{2GM_2}{v_r^2 + c^2_s}, \quad (3)$$

In this expression, $v_r^2 = v_\text{orb}^2 + v_s^2$ is the velocity of wind material near the secondary, $v_\text{orb}$ is the orbital velocity of the secondary, and $c_s$ is the sound speed of the wind. For typical velocities of $\sim 10$–20 km s$^{-1}$, $M_w/M_\text{acc}$ ranges from $\sim 20\%$ for close binaries with $a \approx 3$ AU, to $\lesssim 0.1\%$ for wider binaries with $a \approx 100$ AU. Thus, maximum accretion rates can briefly reach roughly $10^{-5}$ $M_\odot$ yr$^{-1}$ for close binaries with AGB star companions.

Despite the relatively low efficiency, wind accretion can add considerable mass to the secondary. While on the RGB and AGB, the primary star loses most of its initial mass. For primaries with initial masses of $1$–7 $M_\odot$, close companions can accrete as much as $1 M_\odot$ from wind accretion.

Detailed AMR simulations of wind-accretion disks yield results consistent with these simple estimates (de Val-Borro et al. 2009). When the star serves as a sink term in the simulations, circumstellar material forms a fairly stable disk with a total mass of $\sim 10^{-3}$ of the stellar mass out to radii of roughly 5 AU. Typically, $\sim 5\%$ of the mass lost from the AGB star goes through the disk, consistent with the simple Bondi–Hoyle estimates.

To develop a model for the structure of the disk, we assume the primary loses mass at a constant rate, $\dot{M}_w$. For wide binaries, wind material is ejected far from the accretor at approximately parallel trajectories. Gravitational focusing close to the accretor then sets the cross-section for material hitting the disk. Instead of the geometrical cross-section, the fraction of the wind captured by the disk up to the distance $r$ is $\Sigma_{\text{M}_\text{acc}}(<r) \propto r$. For an outer radius equal to $R_a$

$$\Sigma_{\text{M}_\text{acc}}(<R) = \frac{R}{R_a}, \quad (4)$$

The differential accretion rate per radial bin in the disk, is therefore

$$\dot{M}_{\text{acc}}(R) dR = M_{\text{tot}} \frac{dR}{R_a}. \quad (5)$$

The source function, the instantaneous surface density profile of material accreted into the disk, is

$$\Sigma_{\text{f}}(R) dR = \frac{\dot{M}_{\text{acc}}(R)}{2\pi R} = \frac{M_{\text{tot}} \dot{M}_a dR}{2\pi R R_a} \quad (6)$$

As wind material falls into the disk, it dissipates kinetic energy. Together with irradiation from the central star and viscous heating, this energy contributes to the heating of the disk. The amount of energy input into an annulus at radius $r$ is the gravitational binding energy,

$$\dot{E}_r(R) dR = G M_a \dot{M}_{\text{acc}} dR \quad (7)$$

Although the disk does not accrete material beyond $R_a$, viscous processes can expand the disk beyond $R_a$. However, tidal forces from the binary companion limit the outer radius. Typically, this outer radius is roughly 90%-95% of the Roche limit, $a_2 = r_2 a$, where

$$\eta_2 = \frac{0.49 q^{2/3}}{0.6q^{2/3} + \ln(1 + q^{1/3})}, \quad (8)$$

and $q = M_2/M_1$ (Eggleton 1983).

For wind mass loss with $M_1 > M_2$ and low accretion efficiency, binaries expand. We approximate this expansion assuming adiabatic mass loss from the primary star (Hadjidemetriou 1963)

$$a(t) = \frac{M_1 \text{init} + M_2 \text{init}}{M_1(t) + M_2(t)} a_{\text{init}}, \quad (9)$$

where $M_\text{init}$ are the initial masses of the binary components, $a_{\text{init}}$ is the initial separation, and

$$M_i(t) = M_i \text{init} + \int_0^t \dot{M}_i(t) dt, \quad (10)$$

for $i = 1, 2$. Assuming a constant mass-loss rate from the primary, $M_i(t) = M_i \text{init} - M_a t$ and $M_2(t) = M_2 \text{init} + \xi M_a t$, where $\xi$ is the accretion efficiency. For a system losing mass in a wind ($\xi < 1$), $a(t) > a_{\text{init}}$. Thus, the binary separation expands with time. For mass-loss rates and separations of interest, the disk reaches a steady state on timescales shorter than the lifetime of the AGB star (see results in Table 1). Thus, we can safely adopt a simple prescription for the outer radius of the disk

$$R_{\text{out}} = a_2 \quad (11)$$

and assume that $a$ is constant in time. Considering a broad range of $a$ and $M_w$, allows us to derive conditions in the disk for a broad range of plausible binaries (see Table 1).

The inner boundary of the disk, $R_{\text{in}}$, is limited by the radius of the accreting star

$$R_{\text{in}} = R_2. \quad (12)$$

We adopt an inner radius of $1.5 R_2$; the precise value of the inner radius has little impact on the overall structure of the wind-filled accretion disk.

Equipped with the necessary initial and boundary conditions, we now discuss a model for the evolution of the disk.
3. DISK EVOLUTION

3.1. Numerical Calculation

For a disk with surface density $\Sigma$ and viscosity $\nu$, conservation of angular momentum and energy leads to a nonlinear diffusion equation for the time evolution of $\Sigma$ (e.g., Lynden-Bell & Pringle 1974; Pringle 1981),

$$\frac{\partial \Sigma}{\partial t} = 3R^{-1} \frac{\partial}{\partial R} \left( R^{1/2} \frac{\partial}{\partial R} \left( \nu \Sigma R^{1/2} \right) \right) + \left( \frac{\partial \Sigma}{\partial t} \right)_{\text{ext}}. \quad (13)$$

The first term is the change in $\Sigma$ from viscous evolution; the second term is the change in $\Sigma$ from other processes, including mass loss from photoevaporation (e.g., Alexander et al. 2006) or planet formation (e.g., Alexander & Armitage 2009) and mass gain from wind material falling into the disk (given by Equation (6)). Our approach neglects the torque of the disk by the wind, which is small for a disk in a binary system. For the wide binaries we consider (30 and 100 AU separation), the wind momentum can compress the outer edges of the disk but does not affect its bulk structure. The viscosity is $\nu = \alpha c_s H$, where $c_s$ is the sound speed, $H$ is the vertical scale height of the disk, and $\alpha$ is the viscosity parameter. The sound speed is $c_s^2 = \gamma k T_\ast / \mu m_H$, where $\gamma$ is the ratio of specific heats, $k$ is Boltzmann’s constant, $T_\ast$ is the midplane temperature of the disk, $\mu$ is the mean molecular weight, and $m_H$ is the mass of a hydrogen atom. The scale height of the disk is $H = c_s \Omega^{-1}$, where $\Omega = \sqrt{GM_d/R^3}$ is the angular velocity.

To solve this equation numerically, we assume that the midplane temperature is the sum of the energy generated by viscous ($T_d$) dissipation, the energy from irradiation ($T_I$), and the energy of infalling material from the wind ($T_W$),

$$T_d^4 = T_I^4 + T_d^4 + T_W^4. \quad (14)$$

The viscous temperature is

$$T_d^4 = \frac{27 \kappa \nu \Sigma^2 \Omega^2}{64 \sigma}. \quad (15)$$

where $\kappa$ is the opacity and $\sigma$ is the Stefan–Boltzmann constant (e.g., Ruden & Lin 1986; Ruden & Pollack 1991). With $\nu = \alpha c_s^2 \Omega^{-1}$ and $t_2 = (27\sigma/64\alpha) \sim (\kappa \Omega \Sigma^2)$, the viscous temperature is

$$T_d^4 = t_2 T_d. \quad (16)$$

The energy from infalling material follows from Equation (7). Balancing the kinetic energy of infall with the energy radiated by the disk yields

$$\sigma 2\pi Rd RT_w^4(R) = \frac{G M_d \dot{M}_{\text{acc}} dR}{R R_d}. \quad (17)$$

Solving this equation for $T_W$ requires

$$\sigma T_w^4 = \frac{G M_d \dot{M}_{\text{acc}}}{2\pi R^2 R_d}. \quad (18)$$

If the disk is vertically isothermal, the irradiation temperature is $T_I^4(R) = (\theta/2)(R/R_\ast) T_\ast$, where $R_\ast$ and $T_\ast$ are the radius and effective temperature of the central star and (Chiang & Goldreich 1997)

$$\theta = \frac{4}{3\pi} \left( \frac{R_\ast}{R} \right)^3 + \frac{H}{2R} \frac{\partial (H/R)}{\partial R}. \quad (19)$$

Thus, the irradiation temperature is

$$\left( \frac{T_I}{T_\ast} \right)^4 = \frac{2}{3\pi} \left( \frac{R_\ast}{R} \right)^3 + \frac{H}{2R} \left( \frac{R_\ast}{R} \right)^2 \left( \frac{\partial \ln H}{\partial \ln R} - 1 \right). \quad (20)$$
Following Chiang & Youdin (2010) and Hueso & Guillot (2005), we set $\partial \ln H / \partial \ln R = 9/7$. With $H = c_s \Omega^{-1}$, we set $t_0 = (2T_e / 3\pi) (R_e / R)^3$ and $t_1 = (R_e / R)^2 \sim (7 R \Omega)^{-1} \sim (\sqrt{\pi} \mu \mu_m)^{1/2} \sim T_e$. The irradiation temperature is then

$$T_d^4 = t_0 + t_1 t_d^{1/2}.$$  \hfill (21)

Viscous disks are not vertically isothermal (Ruden & Pollack 1991; D’Alessio et al. 1998); however this approach yields a reasonable approximation to the actual disk structure.

Because $T_M$ and $T_d$ are functions of the midplane temperature, and $T_W$ is independent of it, we solve Equation (14) with a Newton–Raphson technique. Using Equations (16), (21), and (18), we rewrite Equation (14) as

$$f(T_d) = T_d^4 - (t_0 + t_1 t_d^{1/2} + t_2 T_d + T_W) = 0.$$  \hfill (22)

Adopting an initial $T_d \approx t_d^{1/3}$ or $T_d \approx t_1^{2/7}$, the derivative

$$\frac{\partial f}{\partial T_d} = 4 T_d^3 - \frac{t_2}{2} T_d^{-1/2} - t_2$$  \hfill (23)

allows us to compute

$$\delta T_d = f \left( \frac{\partial f}{\partial T_d} \right)^{-1}$$  \hfill (24)

and yields a converged $T_d$ to a part in $10^8$ in 2–3 iterations.

In the inner disk, the temperature is often hot enough to vaporize dust grains. To account for the change in opacity, we follow Chambers (2009) and assume

$$\kappa = \kappa_0 \left( \frac{T_d}{T_e} \right)^\alpha$$  \hfill (25)

with $n = -14$ in regions with $T_d > T_e = 1380$ K (Ruden & Pollack 1991; Stepiński 1998). For simplicity, we assume $\kappa = \kappa_0$ when $T_d < T_e$.

To solve for the time evolution of $\Sigma$, we use an explicit technique with $N$ annuli on a grid extending from $x_{in}$ to $x_{out}$ where $x = R^{1/2}$ (Bath & Pringle 1981, 1982a), and $R_{in}$ and $R_{out}$ are inner and outer boundaries determined by Equations (11). To verify this code, Bromley & Kenyon (2011) compare numerical solutions with analytic results from Chambers (2009).

### 3.2. Potential Caveats

Our study explores wind-accretion disks using simple numerical models. These models provide us with the ability to explore the long-term evolution of disks with a wide range of initial conditions. Because these disks reach an approximate steady state, we can test whether the solutions scale as expected with the binary separation and the mass-loss rate of the primary star. However, the approach makes many simplifying assumptions, and does not consider many physical processes known to occur in these systems. Because the physics we do not include is often poorly understood or weakly constrained, it seems prudent to begin with a simple physical structure and add additional physics as needed to understand a broader range of phenomena. The following paragraphs discuss useful physics to consider for future studies.

**Radiative heating of the disk by the companion.** Our calculations do not include radiative heating of the disk by the primary star (mass donor). If the primary has a radius $R_1 \approx 0.5$ AU and a luminosity $L_1$, a flared disk surrounding the secondary with outer radius $R_{out}$ and vertical scale height $h_{out} \approx 0.05 R_{out}$ intercepts 0.4%–0.01% of $L_1$ for $a = 3$–100 AU (Kenyon & Hartmann 1987). With $L_1 \approx 500$–1000 $L_\odot$, this contribution to disk heating is comparable to irradiation by the central MS star. In wide binaries with $a \approx 10$–100 AU, AGB stars can have $R_1 \approx 1$ AU. The disk then intercepts 0.2% ($a = 10$ AU) to 0.02% ($a = 100$ AU) of the radiation from the primary. For typical AGB $L_1 \sim 10^4 L_\odot$ and $a \approx 10$ AU, this component is somewhat larger than irradiation from the central star. In all configurations, the outer rim of the disk intercepts nearly all of this radiation. If the outer disk radiates efficiently, this extra heating should not impact the structure of material in the inner disk. To test this assumption, we performed several test simulations with a fixed outer boundary temperature, $\sigma T_{out}^4 = f_{out} L_1 / 4\pi R_{out} h_{out}$. This extra heating reduces the disk surface density by a factor of $\sim 2$ near the outer boundary but changes the surface density at smaller radii, $R \lesssim 0.95 R_{out}$, by less than 1%. Thus, neglecting this component has little impact on our results. If the outer disk radiates inefficiently, the outer disk will expand and transport thermal energy radially inward. If the scale height of this puff up disk becomes comparable to the accretion radius $R_\star$ it could significantly affect the accretion process, and would require a different approach.

**Disk viscosity.** The processes underlying the origin of the viscosity in accretion disks are still not understood. Here, we follow most studies and adopt a simple prescription using the $\alpha$ viscosity parameter. In particular, we do not directly consider here any interactions of magnetic fields and their evolution, and we use a constant $\alpha$ in each model. Our general results are shown for a choice of $\alpha = 0.01$. Table 1 also provides additional data for the cases of $\alpha = 0.001$ and $\alpha = 0.1$, enabling us to consider the overall dependence of disk structure and evolution on the viscosity parameter.

**Binary orbital expansion.** As briefly discussed above, we do not change the binary separation and mass in our calculation. In most cases, the wind-accretion disk evolves to a steady state in a relatively short timescale. Our approach then provides a good approximation to reality. When evolution in the binary separation occurs more rapidly than the timescale to reach steady-state conditions in the disk (e.g., when the timescale to reach steady state is $\gtrsim 10^3$ yr), our approximation is less accurate. Our results should then be taken with more caution. In general, both the evolution of the binary and the equilibrium disk mass evolve slowly enough to treat the time evolution with a sequence of steady disk models in binaries with increasing semimajor axis.

**Accretion model and disk instabilities.** We use a simplified one-dimensional model for the accretion infall of wind material into the disk. Although this model captures many of the main properties of the material infall, reality is probably more complex. Aside from our failure to treat the complex interaction between the wind and the outer disk, our simplified 1.5D model does not allow for variations in the vertical structure which might lead to thermal instabilities and dwarf nova outbursts (e.g., Cannizzo et al. 2012 and references therein) or for winds driven by radiation pressure from the inner disk (e.g., Noebauer et al. 2010). In particular, the inner region where the temperatures rise above 1000–2000 K (roughly where dust grains evaporate) might become thermally unstable (see, e.g., Alexander et al. 2011). Though the location of the disk instability regime depends on the opacity model, the mass-loss rate of the AGB primary, and other input parameters, Figure 3 shows that...
temperatures as high as 1000 K are typically achieved only for a small region, extending up to an AU for the models with the highest accretion rate and/or the smallest binary separations. In most models, these relatively high temperatures are achieved only over small regions in the inner 0.1–0.3 AU of the disk. In models with the smallest binary separation (3 AU), the outer disk radius is comparable to disks studied in other contexts (e.g., Alexander et al. 2011). In wider binaries with separations of 10–100 AU, the outer disk extends well beyond regions likely capable of sustaining a thermal instability cycle. Thus, potentially unstable regions can extend over a large fraction of the disk in compact binaries (a ≈ 3 AU) but impact a fairly small fraction of the disk in wider binaries (a > 10–100 AU).

We therefore conclude that models exploring small separations, particularly those with the highest accretion rates (e.g., models 1 and 5 in Table 1 below), could be highly susceptible to disk instabilities in a significant fraction of the inner disk. We caution that our steady-state solutions might not well represent these cases. The steady-state solution is likely to well represent likely outcomes in wider binaries and in compact binaries with low accretion rates. In all cases, likely unstable regions are close to the accreting star and could lead to outbursts with distinct observational outcomes. In Alexander et al. (2011), the optical brightness changes little during the outbursts; the bolometric luminosity changes by a relatively small factor of a few for the largest infall rates, though these are larger than those we study.

4. RESULTS

To develop an understanding of the long-term behavior of wind-fed disks in wide binaries, we consider a grid of models in (a, Mwind) space (see Tables 1 and 2 for model parameters and results). For simplicity, we adopt q = 3, M2 = 1 M⊙ (MS secondary), or M2 = 0.6 M⊙ (WD secondary), and assume constant separation and mass-loss rate with time. For each (a, Mwind) pair, we set uwind = 10 km s⁻¹ and cₐ = 20 km s⁻¹, derive Rₙ from Equation (3) and Rout from Equation (11). The disk inner cutoff is Rₙ = 1.4R₂, with R₂ = 1 R⊙ for the MS stars and R₂ = 0.01 R⊙ for the WD stars.

For an initial disk mass Mₐ,0 = 10⁻⁷ M⊙, the initial surface density distribution is Σ = Σ₀R⁻¹e⁻³⁵R/Rₙ with Σ₀ = Mₐ,0/2πRₙRout. Using our numerical solution to the diffusion equation, we continuously add material to the disk inside R₂ and evolve the surface density in time until (and if) the disk reaches a steady state at time τsteady, or up to 1–2 Myr, corresponding to the maximum lifetimes of AGB stars. As long as the steady-state disk mass Msteady is much larger than the initial disk mass, numerical tests show that the initial surface density distribution has little impact on the final steady-state surface density distribution.

We consider the system to reach a steady state when the change in disk mass per unit time is smaller than 10⁻³ of the input rate. In most cases, disks reach steady state on timescales much shorter than the AGB lifetimes (see Figure 1). Several models with the largest binary separations do not properly fulfill this steady-state condition by the end of the simulation, but the actual change in the disk structure at late times is small. For these models we show the disk properties at the end of the simulation run time, τrun. In these cases, τrun is already comparable to or longer than the typical AGB lifetime. Thus, wind-accretion disk cannot achieve a steady-state configuration. Nevertheless, the disk structure in these cases hardly changes over most of the AGB lifetime.

![Figure 1](image-url) Disk mass profile and evolution. Top: the evolution of the total mass in the accretion disk. Disks in binaries with larger separations are wider and accommodate more mass, as they are truncated at distances farther from the star. Bottom: evolution of the accretion rate onto the star. Both upper figures show how the disk evolves into a steady-state configuration over timescales which are typically much shorter than the mass-loss timescales. The lines correspond to the models described in Table 1. Lines of the same color correspond to similar mass loss from the companion (10⁻⁵, 10⁻⁶, 10⁻⁷, and 10⁻⁸ M⊙ yr⁻¹ from top to bottom). Lines of the same type correspond to similar binary separation (3 AU, dash-dotted; 10 AU, dotted; 30 AU, solid; 100 AU, dashed).

Tables 1 and 2 summarize the main results of the calculations. The detailed steady-state profile and evolution of the disks and the accretion rate (on to the star) are shown in Figures 1–3. Although we focus on MS star accretors in the figures, results for WD accretors are similar. Table 2 shows results for WDs with negligible luminosity. If the accreting star is a hot WD with L = 100 L⊙, UV photons heat up disk regions close to the central star but impact the outer disk very weakly. Thus, the time evolution of the accretion rate and disk mass (in Figure 1) and the radial profiles of cumulative mass, surface density, and disk temperature (Figures 2–3) are fairly representative of WD accretors and MS accretors.

For all calculations, the disk mass and the accretion rate onto the star increase roughly linearly with time and then reach an approximate steady state (Figure 1). The timescale to reach steady state increases roughly linearly with increasing binary separation. Although the relation is shallower than a linear one, binaries with smaller mass infall rates also take longer to
reach steady state than binaries with large mass infall rates. The steady-state disk masses range from $\sim10^{-3} M_\odot$ for binaries with massive winds and small separations to $\sim10^{-7} M_\odot$ for binaries with low mass-loss rates and large separations. For models with identical mass-loss rates from the evolved companion, wider binaries have smaller infall rates into the disk and smaller disk masses. When binaries with different $a$ have similar infall rates, wider binaries have larger and more massive disks.

For most MS star models, wind-fed accretion rates produce negligible luminosities. With typical steady accretion rates smaller than $10^{-7} M_\odot$ yr$^{-1}$, the accretion luminosity is comparable to or smaller than the stellar luminosity of $1 L_\odot$. Thus, the accretion disk and boundary layer are invisible (Kenyon & Webbink 1984). For close binaries with AGB-type primary stars, accretion rates of $10^{-6} M_\odot$ yr$^{-1}$ produce modest accretion luminosities of $\sim10 L_\odot$. Although optical radiation from the disk and boundary layer are not detectable, ultraviolet (UV) radiation from the boundary layer is probably visible (Kenyon & Webbink 1984).

Models of wind-fed WD stars are more interesting. In these systems, accretion rates exceeding $10^{-9} M_\odot$ yr$^{-1}$ yield bright UV sources which can ionize the wind from the primary (Kenyon & Webbink 1984). This ionized wind produces bright optical and UV emission lines (Kenyon & Webbink 1984) and a luminous radio sources at cm wavelengths (Taylor & Searquist 1984). These systems would be classified as symbiotic stars (Kenyon 1986). Despite their much lower accretion rates, the central star accretes a larger fraction of infalling material in wide binaries. For $a \approx 3$–100 AU, the Bondi–Hoyle accretion radius $R_a$
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**Figure 3.** Radial temperature profile of wind-accretion disks (inset shows a zoom up region; note linear distance-scale). The radial temperature profiles of wind-accretion disks under various conditions are shown; the lines correspond to similar models as described in Figure 1. Note that regions in which temperatures rise beyond approximately 1000 K might be susceptible to thermal disk instability, not modeled here. (A color version of this figure is available in the online journal.)

exceeds $R_a$ the stagnation radius where the flow velocity in the disk changes sign. Thus, some fraction of infalling material flows out through the disk. The rest accretes onto the central star. In these simulations, the ratio of accreted to ejected material ranges from roughly 55% for close binaries ($a = 3$ AU) to $\sim 80\%$ for wide binaries ($a = 100$ AU). The accretion flow in the accretion disk is bimodal. Generally, material at separations somewhat shorter than $R_a$ flows into the star; the rest of the material outflows into the extended disk up to $R_{\text{out}}$ and is eventually ejected from the disk.

Accreted material typically has a negligible impact on the central MS star. For typical lifetimes of $10^4$ yr or less, the central star accretes less than 1% of its initial mass. In the closest binaries with the highest mass-loss rates, the accretor may accrete from 10% to 90% of its initial mass over the lifetime of the companion (e.g., see the overall accreted mass assuming steady state, over 0.5 Myr, $M_{\text{disk}}$). Although our models do not account for the mass increase of the secondary in these cases, the disk achieves steady state long before the MS star accretes a significant amount of mass. Thus, our steady-state disks still correctly represent the likely physical state of the system.

Accreted material has a more significant impact on accreting WDs. Aside from their visibility as symbiotic stars, wind-fed disks surrounding WDs can produce a variety of eruptive phenomena (Kenyon 1986). At low accretion rates, cold WDs undergo degenerate shell flashes and become classical novae (Kenyon & Truran 1983). WDs accreting at larger rates yield non-degenerate flashes as in the symbiotic stars V1016 Cyg and HM Sge (Mikołajewska & Kenyon 1992). In some systems, the rates are sufficient to allow the WD to evolve close to the Chandrasekhar limit and become a Type Ia supernova (Starrfield et al. 2005).

Steady-state surface density profiles generally follow expectations. Over the entire disk, the surface density approximately scales with radius as $\Sigma \propto r^{-n}$, with $n \approx 0.8-1.5$. The normalization of the surface density scales almost linearly with the accretion rate into the disk. Changes in slope typically occur in the inner disk, when the disk becomes hot enough to evaporate dust grains, and in the outer disk, when the energy generated by irradiation or infall dominates the energy generated by viscosity (e.g., Chambers 2009, and references therein).

We briefly explored the effect of different viscosity parameters on the overall structure and mass of the disks (see Table 1). In general, the steady-state disk mass is roughly inversely proportional to $\alpha$, and scales linearly with $a^{-1}$. These results hold as long as the disk is cold. As the disk gets hotter, more of the disk has an opacity-dependent temperature which slightly changes the relation, but the overall behavior is close to linear. Besides the different overall scaling, the disk spatial structure is very similar.

Steady-state temperature profiles are generally shallower than steady-adequately disks fed by a lobe-filling companion (Lynden-Bell & Pringle 1974). In most steady-state disks, the effective temperature of the disk scales with radius as $T \propto r^{-m}$, with $m \approx 0.75$. For wind-fed disks in very wide binaries at very low accretion rates, the steady-state slope approaches the standard $m \approx 0.75$. As the binary separation contracts and the wind-accretion rate grows, the temperature profile becomes more shallow. For the closest binaries, our results suggest $m \approx 0.4$. The range in the slope is small, from $m \approx 0.4$ in close binaries with large accretion rates, to $m \approx 0.7$ in wide binaries with small accretion rates. Thus, the temperature profile scales very weakly with the separation or the mass-loss rate of the primary star.

5. DISCUSSION

The results from the calculations lead to several broad conclusions. Physical properties of wind-fed accretion disks depend on the orbital separation and the evolutionary state (mass-loss rate) of the primary star. When the secondary is an MS star, it is observable only when the binary is close and the primary is an AGB star. Otherwise, wind-fed accretion has few, if any, direct observational consequences. When the secondary is a WD, the binary is almost always observable as a symbiotic star.

Although optical/UV spectroscopic observations are sufficient to detect luminous symbiotic stars (Kenyon 1986), direct imaging can often reveal the binary companion. In Mira B, IUE spectroscopic observations first identified emission from a hot WD accreting material from the Mira wind (Reimers & Cassatella 1985), recent Hubble Space Telescope observations resolved the system and detected the accretion disk directly (Ireland et al. 2007). For a derived binary separation of $\sim 70$ AU, the inferred accretion luminosity of $\sim 10^{-10} M_\odot \text{yr}^{-1}$ (Sokoloski & Bildsten 2010) is roughly consistent with our models. High-resolution imaging of other symbiotic-like binaries, including R Aqr and CH Cyg, would provide additional tests of these calculations.

Aside from producing interesting symbiotic stars, WD accretors in wide binaries can accrete enough material to approach the Chandrasekhar limit (Table 2; Starrfield et al. 2005; Shen & Bildsten 2007, and references therein). If these WDs grow, they are susceptible to traditional Type Ia supernovae (Arnett 1996) and to sub-Chandrasekhar helium detonation supernovae (e.g., Woosley et al. 1986; Bildsten et al. 2007; Perets et al. 2010; Waldman et al. 2011). In most simulations, the growth of the WD depends on the way material is accreted. Thus, understanding the structure of the disk is important for understanding whether wind-fed WDs are good SNe progenitors. Our disk models provide a step along the path to understanding these outcomes.
Wind-fed disks around WDs provide a less traditional connection to other astrophysical systems. For separations of 3–100 AU and mass-loss rates of \(10^{-8} - 10^{-5} \ M_\odot \, \text{yr}^{-1}\), wind-fed disks have surface density and temperature profiles similar to those observed in low-mass protoplanetary disks. For comparison, Figure 2 shows the range of surface density profiles inferred for an ensemble of protoplanetary disks in nearby star-forming regions (Andrews & Williams 2007; Chiang & Youdin 2010). The shaded region corresponds to disks with masses of \(\sim 2 \times 10^{-4} - 10^{-1} \ M_\odot\). With typical lifetimes of a few Myr, physical processes in protoplanetary disks yield planets on short timescales. Given the similar lifetimes of AGB mass-losing primary stars, it is plausible that some aspects of planet formation occur in the wind-fed disks of wide binary systems (Perets 2010, 2011).

Despite the similarity in instantaneous disk masses, disks in wide binaries have a distinct advantage over protoplanetary disks in producing planets. Continuous feeding by the primary star guarantees that the total mass available for planet formation is, in principle, larger for wide binaries than for protoplanetary disks. In close binaries with evolved AGB primaries, the total mass accreted by the central star is at least as large as most protoplanetary disks. Thus, massive planets could grow in many of these disks.

Testing this idea is challenging. The most massive disks in wide binaries are not as large as those in protoplanetary disks; dust emission from the AGB primary complicates the acquisition of the submillimeter and millimeter observations required to estimate dust masses. Occasional nova eruptions may also frustrate the coagulation processes that grow dust grains from the Mira wind into the planetesimals that produce planets. However, ALMA has sufficient resolution to detect structures in the disks of wide binaries. Identifying structures similar to those in protoplanetary disks might allow robust comparisons between the grain properties and the profiles of surface density and temperature. Detection of debris disks around nearby WDs suggests that planets (or debris from planets) survive the evolution from an MS star into a WD (e.g., Kilic et al. 2012). The frequency of wide WD companions to WD debris disks provides some estimate for the likelihood of this “second-generation” planet formation scenario.

6. SUMMARY

Using a set of numerical calculations, we have explored the evolution and long-term steady-state structure of wind-accretion disks in wide evolved binaries (\(a \approx 3–100 \AU\). These systems evolve rapidly (see Figure 1) and achieve a steady state on timescales which are typically much shorter than typical stellar evolution timescales of \(\sim 10^7\) yr. During steady state, disks have similar surface density and temperature profiles with total masses of a few \(10^{-5} - 10^{-3} \ M_\odot\). The radial density profile is a broken power law in the inner regions and an exponential decline at the outer edge where the disk is truncated by tidal forces from the companion. The radial temperature profile is also described by a broken power law but does not decline dramatically at the outer edge of the disk.

Understanding the formation, structure, and evolution of wind-fed accretion disks is important for a wide variety of phenomena in evolved wide binary systems, including chemically peculiar stars, novae, supernovae, stellar outbursts, and symbiotic binaries. In close binary systems with evolved AGB primaries, significant accretion from a wind-fed disk provides a natural mechanism for chemical enrichment and abundance anomalies in otherwise normal MS and WD stars. Over a broad range of separations and mass-loss rates, wind-fed accretion onto WDs can produce nova eruptions and, possibly, supernova eruptions. By providing a foundation for relating the mass-loss rate of evolved red giants to the temperature and luminosity of the companion, our results also enable more detailed studies of the physical structure of individual binary systems.

Our analysis begins to make a link between the disks in evolved binaries and protoplanetary disks. The structure and evolution of both types of disks depends on uncertain physics, including the disk viscosity and interactions between stellar photons and disk material. For most accretion rates, large disks often drive massive winds which may interact with the surrounding wind from the primary (in a binary system) or a molecular cloud (in a protostellar system). Observational comparisons among disks with similar accretion rates and sizes might improve our overall understanding of accretion phenomena.

REFERENCES

Alexander, R. D., & Armitage, P. J. 2009, ApJ, 704, 989
Alexander, R. D., Clarke, C. J., & Pringle, J. E. 2006, MNRAS, 369, 216
Andrews, R. D., & Wynn, G. A., King, A. R., & Pringle, J. E. 2011, MNRAS, 418, 2576
Andrews, S. M., & Williams, J. P. 2007, ApJ, 671, 1800
Arnett, D. (ed.) 1996, Supernovae and Nucleosynthesis: An Investigation of the History of Matter from the Big Bang to the Present (Princeton, NJ: Princeton Univ. Press)
Bath, G. T., & Pringle, J. E. 1981, MNRAS, 194, 967
Bath, G. T., & Pringle, J. E. 1982a, MNRAS, 199, 267
Bath, G. T., & Pringle, J. E. 1982b, MNRAS, 201, 345
Bildsten, L., Shen, K. J., Weinberg, N. N., & Nelemans, G. 2007, ApJL, 662, 195
Bromley, B. C., & Kenyon, S. J. 2011, ApJ, 735, 29
Busso, M., Gallino, R., Lambert, D. L., Travaglio, C., & Smith, V. V. 2001, ApJ, 557, 802
Cannizzo, J. K., Smale, A. P., Wood, M. A., Still, M. D., & Howell, S. B. 2012, ApJ, 747, 117
Chambers, J. E. 2009, ApJ, 705, 1206
Chiang, E., & Youdin, A. N. 2010, AREPS, 38, 493
Chiang, E. I., & Goldreich, P. 1997, ApJ, 490, 368
D’Alessio, P., Canto, J., Calvet, N., & Lizano, S. 1998, ApJ, 500, 411
de Val-Borro, M., Karovská, M., & Sasselov, D. 2009, ApJ, 700, 1148
Duquennoy, A., & Mayor, M. 1991, A&A, 248, 485
Eggleton, P. P. 1983, ApJ, 268, 368
Hadjidemetriou, J. D. 1963, Icar, 2, 440
Han, Z., Eggleton, P. P., Podsiadlowski, P., & Tout, C. A. 1995, MNRAS, 277, 1443
Hueso, R., & Guillot, T. 2005, A&A, 442, 703
Ireland, M. J., Monnier, J. D., Tuthill, P. G., et al. 2007, ApJ, 662, 651
Kenyon, S. J. 1986, The Symbiotic Stars (Cambridge: Cambridge Univ. Press), 295
Kenyon, S. J., & Hartmann, L. 1987, ApJ, 323, 714
Kenyon, S. J., Oliverse, N. A., Mikolajewska, J., et al. 1991, AJ, 101, 637
Kenyon, S. J., & Truran, J. W. 1983, ApJ, 273, 280
Kenyon, S. J., & Webbink, R. F. 1984, ApJ, 279, 252
Kenyon, S. J., Webbink, R. F., Gallagher, J. J., & Truran, J. W. 1982, A&A, 106, 109
Kilic, M., Patterson, A. J., Barber, S., Leggett, S. K., & Dufour, P. 2012, MNRAS, 419, L59
Livio, M., Soker, N., de Kool, M., & Savonije, G. J. 1986, MNRAS, 222, 235
Luck, R. E., & Bond, H. E. 1991, ApJS, 77, 515
Lynden-Bell, D., & Pringle, J. E. 1974, MNRAS, 168, 603
Mastrodemos, N., & Morris, M. 1998, ApJ, 497, 303
McClure, R. D., & Woodsworth, A. W. 1990, ApJ, 352, 709
Mikolajewska, J., & Kenyon, S. J. 1992, MNRAS, 256, 177
Noebauer, U. M., Long, K. S., Sim, S. A., & Knigge, C. 2010, ApJ, 719, 1932
Olofsson, H., González Delgado, D., Kerschbaum, F., & Schöier, F. L. 2002, A&A, 391, 1053
Perets, H. B. 2011, AIP Conf. Proc. 1331, Planetary Systems beyond the Main Sequence (Melville, NY: AIP), 56
Perets, H. B. 2010, arXiv:1001.0581
Perets, H. B., Gal-Yam, A., Mazzali, P. A., et al. 2010, Natur, 465, 322
Pringle, J. E. 1981, ARA&A, 19, 137
Ramstedt, S., Schöier, F. L., Olofsson, H., & Lundgren, A. A. 2006, A&A, 454, L103
Reimers, D., & Cassatella, A. 1985, ApJ, 297, 275
Ruden, S. P., & Lin, D. N. C. 1986, ApJ, 308, 883
Ruden, S. P., & Pollack, J. B. 1991, ApJ, 375, 740
Shen, K. J., & Bildsten, L. 2007, ApJ, 660, 1444
Siviero, A., Munari, U., Dallaporta, S., et al. 2009, MNRAS, 399, 2139
Soker, N., & Rappaport, S. 2000, ApJ, 538, 241
Sokoloski, J. L., & Bildsten, L. 2010, ApJ, 723, 1188

Sokoloski, J. L., Luna, G. J. M., Mukai, K., & Kenyon, S. J. 2006, Natur, 442, 276
Starrfield, S., Hix, W. R., Timmes, F. X., et al. 2005, NuPhA, 758, 455
Stepinski, T. F. 1998, Icar, 132, 100
Taylor, A. R., & Seaquist, E. R. 1984, ApJ, 286, 263
Waldman, R., Sauer, D., Livne, E., et al. 2011, ApJ, 738, 21
Wang, Y.-M. 1981, A&A, 102, 36
Webbink, R. F., Livio, M., Truran, J. W., & Orio, M. 1987, ApJ, 314, 653
Woosley, S. E., Taam, R. E., & Weaver, T. A. 1986, ApJ, 301, 601