ROTATING STARS AND THE FORMATION OF BIPOLAR PLANETARY NEBULAE. II. TIDAL SPIN-UP

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ABSTRACT

We present new binary stellar evolution models that include the effects of tidal forces, rotation, and magnetic torques with the goal of testing planetary nebulae (PNs) shaping via binary interaction. We explore whether tidal interaction with a companion can spin-up the asymptotic giant branch (AGB) envelope. To do so, we have selected binary systems with main-sequence masses of 2.5 Mₜₜ and 0.8 Mₜₚ and evolve them allowing initial separations of 5, 6, 7, and 8 au. The binary stellar evolution models have been computed all the way to the PNs formation phase or until Roche lobe overflow (RLOF) is reached, whatever happens first. We show that with initial separations of 7 and 8 au, the binary avoids entering into RLOF, and the AGB star reaches moderate rotational velocities at the surface (∼3.5 and ∼2 km s⁻¹, respectively) during the inter-pulse phases, but after the thermal pulses it drops to a final rotational velocity of only ∼0.03 km s⁻¹. For the closest binary separations explored, 5 and 6 au, the AGB star reaches rotational velocities of ∼6 and ∼4 km s⁻¹, respectively, when the RLOF is initiated. We conclude that the detached binary models that avoid entering the RLOF phase during the AGB will not shape bipolar PNs, since the acquired angular momentum is lost via the wind during the last two thermal pulses. This study rules out tidal spin-up in non-contact binaries as a sufficient condition to form bipolar PNs.

Key words: binaries: general – planetary nebulae: general – stars: AGB and post-AGB – stars: evolution – stars: magnetic fields – stars: rotation

1. INTRODUCTION

During nearly 30 years of research in the formation of planetary nebulae (PNs), several explanations have been proposed in the literature to produce the bipolar shapes. The most popular explanation involves the presence of a dense medium in the equatorial latitude which collimates the later fast stellar wind (Calvet & Peimbert 1983; Balick 1987; Icke 1988; Icke et al. 1989; Mellema et al. 1991; Frank & Mellema 1994). A number of models focus on how to form a dense medium along the equator of the asymptotic giant branch (AGB) star through stellar rotation (Ignace et al. 1996; García-Segura 1997; García-Segura et al. 1999), while others concentrate on magnetic torques, succeeding only sometimes at creating such a dense medium (Matt et al. 2000).

A long debate regarding the need for a binary companion has also taken place during the past three decades, since AGB wind asphericities could naturally arise from the interaction of AGB stars with a companion in several ways. The most popular involve common envelope evolution and the initial phase of spiral-in (Livio 1993; Soker 1997; Sandquist et al. 1998; Nordhaus & Blackman 2006; De Marco 2009; Ricker & Taam 2012; Passy et al. 2012; De Marco et al. 2013), gravitational focusing (Gawryszczak et al. 2002), wind-capture disks (Huarte-Espinosa et al. 2013), and Wind-Roche-Lobe-Overflows (Podsiadlowski & Mohamed 2007; Mohamed & Podsiadlowski 2012). Along these lines, in our paper I (García-Segura et al. 2014), we provided a solid study showing that bipolar PNs formation cannot result from the evolution of a single AGB star, given that these stars do not carry the necessary angular momentum, or to be more specific, the necessary surface rotation.

In this second paper, we explore the alternative of binary (versus single) evolution for bipolar PNs formation from the view point of stellar evolution models. We calculate the effects that a secondary star has on the primary stellar surface and see if these effects can explain the formation of a bipolar nebula. To do so, we compute the first binary models of AGB stars allowed to be spun-up by tidal interactions, and study which rotational velocity we can obtain using this approach. As we mentioned in paper I, surface rotational velocity values above 2 km s⁻¹ could in principle produce considerable asymmetries in these stars.

The spin-up of stellar envelopes has only been calculated for massive stars (see Detmers et al. 2008) to explain the formation of gamma-ray bursts. However, we recall that although the present calculations are the first in the context of the formation of bipolar PNs, the idea has been around for quite some time. Livio (1994) first mentioned tidal spin-up as a possible origin for bipolar PNs, and Soker (1998) later suggested that as a consequence of the spin-up, the primary’s envelope could rotate at several percent of the breakup velocity, which would probably result in a small, but non-negligible, effect on the mass-loss geometry. In this paper, we precisely focus on the above ideas and compute the importance of spin-up in this context. The outline of the paper is as follows. Section 2 describes the numerical scheme and physical approximations, as well as the inputs in our calculations. Section 3 shows the results of the stellar evolution calculations. Finally, in Section 4, we discuss the results and provide the main conclusions of this study.
2. STELLAR MODELS, METHODS, AND PHYSICAL ASSUMPTIONS

The stellar evolution calculations have been performed using the Binary Evolution Code (BEC; Petrovic et al. 2005; Yoon et al. 2006). BEC is a one-dimensional, hydrodynamic, stellar evolution code designed to evolve stellar models of single and binary stars, which originates from the STERN code (Langer 1991). The structure and evolution of stars is governed by a set of partial differential equations, that is, the so-called stellar structure equations. In BEC, which uses the hydrodynamic version, the inertia terms are used (see Kozyreva et al. 2014). This code includes diffusive mixing due to convection, semi-convection (Langer et al. 1985), and thermohaline mixing, as in Wellstein et al. (2001) and Cantiello & Langer (2010).

In rotating stars, centrifugal forces act on the stellar gas leading to deviations from spherical symmetry. For slow to moderate rotation, these deformations remain rotationally symmetric (Tassoul 1978) and no triaxial deformations are expected. The shapes of the equipotential surfaces are affected by the centrifugal potential, and therefore deviate from spherical symmetry. In this scenario, the momentum equation and the energy transport equation for spherically symmetric stars have to be modified. To take this effect into account, the effect of the centrifugal force on the stellar structure is calculated following the method of Kippenhahn & Thomas (1970) in the approximation of Endal & Sofia (1976), and applied to the hydrodynamic stellar structure equations (Heger et al. 2000). Thus, instead of spherical shells, the mass shells correspond to surfaces of constant pressure. The corrections are applied to the acceleration and the radiative temperature gradient. According to Zahn (1975), Chaboyer & Zahn (1992), and Zahn (1992), anisotropic turbulence acts much more strongly on isobars than in perpendicular directions. This enforces shellular rotation rather than cylindrical rotation, and it removes compositional differences on isobaric surfaces. Therefore, it can be assumed that the matter on such surfaces is chemically homogeneous. Together with the shellular rotation, this allows us to retain the one-dimensional approximation (Heger et al. 2000).

Time-dependent chemical mixing and transport of angular momentum due to rotationally induced instabilities, including the shear instability and the Golreich–Schubert–Fricke instability, and the Eddington–Sweat circulations are also included as a diffusive process (Heger et al. 2000). We also include the transport of angular momentum due to magnetic fields (Spruit 2002), as in Heger et al. (2005) and Petrovic et al. (2005). They considered dynamo action in radiative layers, resulting from winding up magnetic fields due to differential rotation, which creates toroidal fields, and the generation of poloidal fields from toroidal fields due to the Tayler instability (the Tayler-Spruit dynamo). The strength of the magnetic fields is calculated using the steady state solution of the Tayler-Spruit dynamo, which is given as a function of the degree of differential rotation. The magnetic field therefore is instantaneous and has no memory of previous timesteps. In our cases, the strongest magnetic field is found in the boundary layer between the core and the envelope where the degree of differential rotation is largest, while in the rest of the envelope, magnetic fields are very weak. The rotation profile in the convection zone is self-consistently calculated by solving the diffusion equation for the transport of angular momentum.

The physics of binary interactions was implemented by Braun (1997) and Wellstein (2001). We assume a circular orbit and adopt the Eggelton’s approximation for the Roche lobe radius (RLOF; Eggleton 1983) for our considered binary systems. We calculate the mass transfer rate due to RLOF following the method by Ritter (1988) in an implicit manner (note, however, that our models do not involve mass transfer in this article). We adopt the prescription by Podsiadlowski et al. (1992) to calculate the evolution of the binary orbit resulting from mass transfer and stellar winds mass loss, and the numerical result by Brookshaw & Tavani (1993) to calculate the amount of specific angular momentum of the stellar wind material that escapes the binary system.

The synchronization of the binary system is computed with the same method as in Wellstein et al. (2001), using the synchronization time given by Zahn (1977) for convective stars:

$$\tau_{\text{sync}} \propto q^{-2} \left( \frac{d}{R} \right)^{6},$$

(1)

where \( q = M_2/M_1 \) is the mass ratio of the binary, with \( M_1 \) being the mass of the primary star that evolves up to the AGB phase and \( M_2 \) the mass of the secondary main-sequence star, \( d \) is the orbital separation, and \( R \) the radius of the primary AGB star. The angular momentum exchange \( \Delta J \) that is added to the whole AGB star at each time step \( \Delta t \) due to tides is

$$\Delta J = (J_{\text{AGB}} - J_{\text{sync}})(1 - e^{-\Delta t/\tau_{\text{sync}}}),$$

(2)

where \( J_{\text{sync}} \) is the spin angular momentum that the AGB star would have when the system is in synchronous rotation.

Recently, Paxton et al. (2015) implemented binary physics into the MESA code by including many ingredients from BEC. Since the two codes use very similar physics, a comparison of the two would not be very meaningful (see Section 2.6 for Numerical Tests in Paxton et al. 2015). There exists no other code that includes both binarity and differential rotation at the present. The accuracy of our results is more dependant on the synchronization physics rather than on numerics because the numerical implementation is quite straightforward. Along these lines, there are two articles which are useful to mention. In the first, Khalilulin & Khalilullina (2010) found that Zahn’s theory is consistent with observations of 101 eclipsing binaries with an early-type, main-sequence component, and in the second one, Zamanov et al. (2007) report evidence for the synchronization of 29 S-type symbiotic stars, which are similar types of binaries. The assumption of circular orbits is a good approximation in our case, since the circularization timescale is comparable to the synchronization timescale. In the case that the orbits were very eccentric, there might be episodes of mass transfer when the stellar components become close enough, for a given mean orbital separation. Considering such a situation would be very interesting, but will be the subject of future work (see Staff et al. 2015 for a recent work along these lines).

The majority of Solar-type stars did seem to have stellar companions (Abt & Levy 1976; Duquennoy & Mayor 1991). The fraction of main-sequence binaries was found to be in the range from 65% to 100% with only about 30% of G-dwarf primaries having no companion up to about 0.01\( M_\odot \) (Duquennoy & Mayor 1991). The multiplicity fraction, however, seems to have a dependence on the mass of the primary with binary systems being about 50% between spectral...
The Astrophysical Journal, 823:142 (8pp), 2016 June 1

GARCÍA-SEGURA ET AL.

3. RESULTS

The evolution of the surface equatorial rotational velocity during the thermal pulsing AGB phase is shown in Figure 1 for each of the initial binary parameters listed in Table 1. We do not show the earlier evolution because the tidal spin-up is not important since the synchronization time is quite dependent on the radius (see Equation (1)). Only when the stellar radius during the AGB is large enough for a given binary separation do tidal forces become important and the synchronization time becomes short enough that the spin-up becomes evident. For comparison, Figure 1 also shows (dotted line) the results of the model presented in paper I in which we included angular momentum transport induced by magnetic torques inside the star.

In Figure 1, we see that all of the binary models have larger surface rotational velocities than the single stellar models computed in paper I (where surface rotational velocities larger than $1 \text{ km s}^{-1}$ were never reached at the late thermal pulsing AGB phase). We find that the models with initial separations of 5 and 6 au end up in RLOF phases while the models with initial separations of 7 and 8 au avoid RLOF altogether. The AGB stars, in the more separated binary models, reach moderate rotational velocities at the surface ($\sim 3.5$ and $\sim 2 \text{ km s}^{-1}$ for models with 7 and 8 au, respectively) during the interpulse phase, just before the last two thermal pulses, and for the closest binary separations explored, 5 and 6 au, the AGB stars reach rotational velocities of $\sim 6$ and $\sim 4 \text{ km s}^{-1}$, respectively, when the RLOF is initiated. Note that the models that do not enter into the RLOF phase lose most of their gained angular momentum via the wind during the last two thermal pulses. This is clearly reflected in Figure 1 with the quick drop of velocity at the end of their AGB evolution.

We find that the maximum rotation velocity gained at the thermal pulsing AGB phase does not depend on the initial velocity assumed for the primary, but only on the initial separation of the binary. To illustrate this, we plot the spin angular momentum of the primary star in Figure 2. We find that the final spin angular momentum reached by the star is independent of the initial assumption, since the major contribution to the spin angular momentum of the primary comes from the orbital angular momentum. However, there is an important transfer of orbital angular momentum to the primary spin angular momentum, the separation (and the orbital period) growth in time due to mass loss, as can be seen in Figure 3. This figure also shows that models c5 and r5 have almost reach corotation at the time our computation is stopped (because they reach RLOF). Models c6 and r6 are slightly below the corotation speed.

One of the most important results of the comparison allowed by our models is that those binaries that do not reach RLOF lose their gained angular momentum through mass loss at the end of the AGB phase. This means that for these models, the formation of bipolar nebulae via magnetohydrodynamical collimation is not possible, since at the time when a PN is formed, after an uncertain amount of transition time, the stellar surface rotation is very small and the effect that the secondary might have had in speeding up the primary will be negligible due to the large separation. Since the sound speed at the surface of an AGB star is of the order of $\approx 1 \text{ km s}^{-1}$, values below this

| Model | $d$ (au) | $M_1 (M_\odot)$ | $M_2 (M_\odot)$ | $v_{\text{eq}}$ (km s$^{-1}$) | Description |
|-------|---------|-----------------|-----------------|-----------------|-------------|
| c5    | 5       | 2.5             | 0.8             | $\sim 0$       | corotation  |
| c6    | 6       | 2.5             | 0.8             | $\sim 0$       | corotation  |
| c7    | 7       | 2.5             | 0.8             | $\sim 0$       | corotation  |
| c8    | 8       | 2.5             | 0.8             | $\sim 0$       | corotation  |
| r5    | 5       | 2.5             | 0.8             | 250            | fast rotator |
| r6    | 6       | 2.5             | 0.8             | 250            | fast rotator |
| r7    | 7       | 2.5             | 0.8             | 250            | fast rotator |
| r8    | 8       | 2.5             | 0.8             | 250            | fast rotator |

types B4 and A7 (Kouwenhoven et al. 2007) and about 33% for F6-K3 stars using the more recent estimates by Raghavan et al. (2010). These values are consistent with the fraction of binary systems found in central stars of PNs (Douchin et al. 2015). Note that both the binary period distribution and the binary fraction are meaningful in performing this comparison.

In the current study, we use binary systems where the primary stars have an initial mass of $2.5M_\odot$. The main reason that motivated the choice of the initial progenitor mass is the higher-mass stellar progenitor that the bipolar PNs morphological class seem to have, as indicated by their chemical abundances (see, e.g., Manchado 2000; Stanghellini et al. 2006), and their closer distribution to the galactic plane (Calvet & Peimbert 1983; Corradi & Schwarz 1995; Manchado 2000; see also paper I Garcia-Segura et al. 2014). For the binary system, we have chosen as the secondary a stellar companion with $0.8M_\odot$. These stars represent some of the most numerous stars on the main sequence and have also been found to be numerous as secondaries in the study of Sco OB2 by Kouwenhoven et al. (2005). A mass fraction $q = M_2/M_1$ of 0.32 seems to be reasonably likely and consistent with the values observed by Kouwenhoven et al. (2005). Note that Duquennoy & Mayor (1991) also found that the binary mass distribution peaks around $q = 0.3$ for Solar-type stars, concluding that, in general, binaries can be formed by an association of stars built from the same IMF. The secondary star is treated as a point mass in this study.

For the stellar mass-loss rate, we adopted Reimers (1975) with $\eta = 0.5$ during the Red Giant phase and the Vassiliadis & Wood (1993) parameterization for the AGB phase.

We have computed two sets of models (eight in total) from the zero age main sequence (ZAMS) to the AGB phase using primaries with two different initial stellar rotation rates and four different binary initial orbital separations. The binary parameters of all of the computed models at the ZAMS are listed in Table 1. In the first set of models, we consider a very small initial rotation, nearly zero, which is the rotation that would result from synchronization at the ZAMS. For the second set of models, we have adopted an initial equatorial rotation velocity of $250 \text{ km s}^{-1}$ at the ZAMS representative of single stars (Fukuda 1982) with the mass of the primary. This velocity has also been chosen to allow for comparisons with the results from paper I. The initial orbital separations used in the calculations are 5, 6, 7, and 8 au. To make the choice for the separations of the binaries, we first computed preliminary basic models that allowed us to select two cases with RLOF and two cases that avoid RLOF, as will be discussed in the next section. The minimum binary separation of 5 au has been tuned to ensure that binary interaction is achieved during the late AGB phase, avoiding altogether a strong tidal interaction during both the main sequence and the Red Giant Branch phase.
value for the rotation will not have any impact on the wind-compression mechanism (Bjorkman & Cassinelli 1993; Ignace et al. 1996). When the calculations are stopped, we obtain surface rotational speeds of 0.03, 0.03, 0.4, and 0.09 km s\(^{-1}\) for models c7, c8, r7, and r8, respectively. The stellar parameters are shown in Table 2.

On the other hand, the models that reach RLOF, generate, for a timespan that requires further calculation, the necessary dynamical conditions to form an equatorial density enhancement in the circumbinary medium, since the rotation is fast (~6 km s\(^{-1}\)) and most likely sustained for the necessary time. These models are promising in forming bipolar nebulae, as we will discuss in the next section. Table 3 gives the parameters of the binary systems approaching the RLOF phases at the time the computations are stopped.

4. DISCUSSION AND CONCLUSIONS

Our set of binary calculations shows that it seems to be a very reasonable assumption that those models that avoid RLOF phases will not form bipolar PNs, since all of the gained

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**Figure 1.** Evolution of the surface equatorial rotational velocity during the thermal pulsing AGB phase. Left: models with ZAMS velocities nearly zero, c5, c6, c7, and c8. Right: models with ZAMS velocities of 250 km s\(^{-1}\), r5, r6, r7, and r8. The initial binary separation is labeled in the upper left corner. The dotted lines correspond to the magnetic model of paper I as a reference.
angular momentum during the AGB phase will be lost via the wind in the last two thermal pulses. It is also important to note the fact that the whole timescale affected by this process is very long, of the order of 100,000 years, taking into account the last thermal pulse and the possibility of a few thousand years more of uncertain transition time until the central star is hot enough for the fast wind mechanism to operate and shape the nebula (see e.g., Villaver et al. 2002). Any fingerprint of a transient, equatorial density enhancement will be dissipated at the time of PN formation in these models. Then, only those models that reach their RLOF phases are promising to explore bipolar PNs shaping, since their evolution will proceed differently than those for single stars or non-contact binaries.

Following our results, it will be extremely useful to have the distribution and latitudinal dependence of the stellar wind coming out from the binary system, in order to model the formation of a bipolar PNs. Unfortunately, the equations by Bjorkman & Cassinelli (1993) and Ignace et al. (1996), which study the formation of wind-compression zones and wind-compression disks around single stars, cannot be applied directly in this case because the geometry is different for binary systems and the center of mass of the binary is not at the center of the AGB star. Currently, there is no analytical or semi-

Figure 2. Same as Figure 1, except for the evolution of the spin angular momentum.
analytical study that predict the distribution of gas in the nearby circumbinary medium for these binary stars, except for the studies where the outflow comes from the Lagrange point $L_2$ (Shu et al. 1979), as we will discuss below.

We have performed a rough estimate of how important a wind-compression mechanism could be by using the results of our models and keeping in mind the applicability restrictions outlined above. We focus the discussion on the last output from model c5. In this model, the AGB star is at 19% of its critical rotation speed which, according to Ignace et al. (1996), can be translated into a density contrast ratio equator/pole of $\sim 9$.

Figure 3. Evolution of the spin period (solid line) and orbital period (dotted line) during the thermal pulsing AGB phase for models c5, c6, c7, and c8 (left panels) and models r5, r6, r7, and r8 (right).

### Table 2

| Model | $M (M_{\odot})$ | $\log L (L_{\odot})$ | $T_{\text{eff}}$ | $R (R_{\odot})$ | $v_{\text{rot}}$ (km s$^{-1}$) | $P_{\text{orbital}}$ (days) |
|-------|----------------|----------------------|-----------------|----------------|-----------------|---------------------|
| c7    | 0.819          | 4.134                | 2970            | 441            | 0.03            | 11,949               |
| c8    | 0.836          | 4.188                | 3148            | 417            | 0.03            | 15,164               |
| r7    | 0.841          | 4.212                | 3164            | 424            | 0.44            | 11,601               |
| r8    | 0.816          | 4.265                | 3159            | 452            | 0.09            | 15,631               |
that these numbers are given in the non-inertial reference frame of the AGB star but we can translate them to the inertial reference frame at the center of mass. The total mass of the binary system is $3.006\, M_\odot$ (accounting for the mass lost) and the center of mass is located $306\, R_\odot$ from the center of the AGB star, i.e., $\sim 26\%$ inside of its envelope. With this in mind, the equatorial surface at the opposite side of the secondary has a velocity of $14\, \text{km s}^{-1}$, which translates to close to $\sim 50\%$ of the critical rotation at that location. We see that $50\%$ of the critical rotation is within the range of conditions for the wind-compression disk formation, of course, if the equations of Ignace et al. (1996) could be applied in the scenario we have just described. What this exercise emphasizes is that mass loss is expected to be important on the side opposite the secondary, in the direction of the Lagrange point $L_3$, where the escape velocity at the surface is smaller than the value that an isolated star would have. This mass-loss enhancement is expected to occur with the shape of a spiral arm.

It is also important to mention that the system will lose an important amount of mass through the Lagrange point $L_2$. This mass loss is expected to have the shape of a spiral arm as well, since the mass ratio $q = 0.36$ at the beginning of the RLOF phase is inside the range $0.064 \leq q \leq 0.78$ (see, e.g., Shu et al. 1979; Mohamed & Podsiadlowski 2012; Pejcha et al. 2015). It is also interesting to note in this context that the ratio $q$ increases along the subsequent evolution as the primary loses mass and the secondary accretes part of it. Once the ratio $q$ exceeds the value of 0.78, if the system is still in RLOF, the ejection of gas through $L_2$ piles up at a finite radius and will form a bounded mass-loss ring instead of an unbound outflow (Shu et al. 1979).

How much time the system will remain in the RLOF phase depends on the rate at which the AGB radius expands, how soon the common envelope is formed, the rate at which orbital angular momentum is lost (shrinking the orbit of the binary system), and on the mass-loss rate from the binary system (increasing the orbital distance). In any case, the proper computations of the mass loss through the Langrangians points $L_2$ and $L_3$ is challenging. These new scenarios are of great interest for future hydrodynamical computations for the formation of bipolar PNs such as those recently carried out by Staff et al. (2015).

Based on our binary stellar evolutionary models with initial $q = 0.32$, we find that the maximum rotation velocity during the AGB phase, as a result of tidal spin-up, is independent of the initial rotational velocity at the ZAMS and only depends on the initial separation.

We find it to be very unlikely that binaries which avoid the RLOF phases could form bipolar PNs with an equatorial waist. The angular momentum lost by the stellar wind in the last thermal pulses, and the unknown transition time from the AGB to the PN formation, will dissipate any transient, equatorial density enhancement formed in the nearby circumbinary medium. The same reasoning applies to the engulfing of giant planets if they do not produce the ejection of the envelope, since no matter how much angular momentum is gained by the AGB star, the last thermal pulses will carry it away via the stellar winds.

The above discussion is only focussed on the classical formation of bipolar nebulae, in which a slow, equatorially denser wind inhibits the future expansion of a fast wind in the equatorial regions. However, this is not the only model that accounts for the formation of bipolarity, since jets and collimated outflows may form around the companion (Soker & Rappaport 2000; Huarte-Espinosa et al. 2013) and power PNs, as they are known to do in Symbiotics (Corradi & Schwarz 1993). These scenarios are based on the formation of an accretion disk around the secondary star where the ejected collimated gas forms the bipolar nebula before the formation of a fast wind by the central star. The separation of the binary system is crucial here in determining the energy of the outflow. According to the separation of the binary, the systems can be classified into three types: Wind-Capture disks for large separations; Wind-RLOF disks for small separations, but still larger than the distances required for RLOF; and the RLOF disks formed at the RLOF phases. Note that systems with small separations (such as models c5, c6, r5, and r6) will pass through the three different types consecutively as the AGB radius increases during stellar evolution. In the Wind-Capture disk, a companion orbiting through the dense wind of an AGB star will capture enough wind material to form an accretion disk and (most likely) power a jet (Huarte-Espinosa et al. 2013). The Wind-RLOF disk is a relatively new mechanism which allows far larger orbital separations than in traditional RLOF. High accretion rates are possible with Wind-RLOF scenarios and the determination of its presence requires solutions for the AGB wind structure (Podsiadlowski & Mohaded 2007; Mohamed & Podsiadlowski 2012). Finally, we have the outflows from RLOF phases where the amount of accretion is largest. Although neither of these possibilities are accounted for in the present study, they are not exclusive, in the sense that the three types, especially the last two (Wind-RLOF and RLOF), will also form a density enhancement in the orbital plane with a spiral shape, due to the mass loss through the Lagrange point $L_2$. In other words, the accretion disks will form both jets and density enhancement in the orbital plane.

Since the orbital separation increases very quickly at the last thermal pulses due to the large mass loss (Figure 3), the wind accretion is expected to decay with time in the first two scenarios. Thus, it is very likely that the formation of bipolar PNs follows the evolution of binary systems with separations where RLOF phases are attained, as well as the subsequent (or imminent) common envelope phases. Proper computations of the time spent in each of the three phases will be extremely important and will be the subject of future studies.

Finally, another point which needs to be considered in this study is the role of magnetic fields. The creation of fields by companions in common envelope scenarios has been discussed.

Table 3
Primary Stellar Parameters when Roche Lobe Overflow is Approached

| Model | \(M (M_\odot)\) | \(\log L (L_\odot)\) | \(T_{\text{eff}}\) | \(R (R_\odot)\) | \(v_{\text{esc}}\) (\(\text{km s}^{-1}\)) | \(d (R_\odot)\) | \(P_{\text{orbital}}\) (days) |
|-------|----------------|----------------|----------------|----------------|----------------|----------------|----------------|
| c5    | 2.206          | 4.147          | 3076           | 415            | 6.1            | 1151.53        | 2612.62        |
| c6    | 1.869          | 4.226          | 2887           | 518            | 4.9            | 1480.16        | 4040.97        |
| r5    | 2.239          | 4.149          | 3086           | 415            | 5.8            | 1144.73        | 2575.88        |
| r6    | 1.999          | 4.239          | 2923           | 513            | 4.8            | 1443.27        | 3799.25        |
in the literature and offers another route whereby binary stars produce bipolar PNs (Tout & Regós 2003; Nordhaus & Blackman 2006). Note that the spin-up of the envelope in the above studies is the direct product of a spiral-in process during the common envelope evolution. In our cases, however, the spin-up of the envelope is due to tidal forces. In the case of non-contact binaries, we found that magnetic fields are very weak inside of the envelope where the degree of differential rotation is negligible. The rotation velocity become too small at the end of the AGB phase and would not display any big difference from single stars. We find it to be very unlikely that magnetic fields could be of any importance here, and they will probably only have a role in forming cool spots (Frank 1995; Soker 2001).

On the other hand, for the closest systems, it would make a great difference. Since AGB stars rotate at $\sim$20$^-$$^\sim$50$^\circ$ of their critical rotation, this is significant enough to make an alpha-omega dynamo efficient. However, as we discussed above, they will become a common envelope soon once mass transfer starts, and the later evolution would be governed by the common envelope. So, most likely, even in this case, the role of a convective dynamo might be minor, although this requires further investigation, since the accumulated magnetic energy could be important for the ejection of the envelope.

In conclusion, this study rules out tidal spin-up in non-contact binaries being a sufficient condition to form bipolar PNs.

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