COLLIMATED OUTFLOW FORMATION VIA BINARY STARS: THREE-DIMENSIONAL SIMULATIONS OF ASYMPTOTIC GIANT BRANCH WIND AND DISK WIND INTERACTIONS

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ABSTRACT

We present three-dimensional hydrodynamic simulations of the interaction of a slow wind from an asymptotic giant branch (AGB) star and a jet blown by an orbiting companion. The jet or “collimated fast wind” is assumed to originate from an accretion disk that forms via Bondi accretion of the AGB wind or Roche lobe overflow. We present two distinct regimes in the wind-jet interaction determined by the ratio of the AGB wind to jet momentum flux. Our results show that when the wind momentum flux overwhelms the flux in the jet, a more disordered outflow results with the jet assuming a corkscrew pattern and multiple shock structures driven into the AGB wind. In the opposite regime, the jet dominates and will drive a highly collimated, narrow-waisted outflow. We compare our results with scenarios described by Soker & Rappaport and extrapolate to the structures observed in planetary nebulae (PNs) and symbiotic stars.

Subject headings: binaries: close — planetary nebulae: general — stars: AGB and post-AGB

1. INTRODUCTION

The role of binary stars in the shaping of planetary and proto-planetary nebulae (PNs, pPNs) has been a topic of debate for many years (Bujarrabal et al. 2000). PNs occur in a wide variety of shapes ranging from spherical to elliptical, narrow-waisted bipolar to bilobed (an elliptical core with projecting bipolar lobes). Some PNs also show narrow jetlike features. pPNs show striking bipolar or jetlike morphologies. PPBs show point-symmetric morphologies in which nebular features are reflected about the central star (Sahai & Trauger 1998). Such features are difficult to explain with purely hydrodynamic models. pPNs have included, either implicitly or explicitly, the presence of a binary companion (see Balick & Frank 2002 for a review of PNs shapes and shaping mechanisms).

Hydrodynamic generalized interacting stellar wind (GISW) models in which successive wind phases from a single evolving asymptotic giant branch (AGB) star sculpt the nebula were believed for some time to hold considerable promise (Balick 1987; Icke, Preston, & Balick 1989). In these scenarios, fast winds from the PNs central star expand into an aspherical slow wind expelled by the AGB progenitor. These models did not explicitly require the presence of a companion, but it was often assumed that the shaping of the slow wind occurred via tidal or common envelope interactions (Livio 1993).

While hydrodynamic models were successful at recovering many elliptical and bilobed nebula, simulations demonstrated that the more extreme “wasp-waisted” bipolars could not be easily recovered (Icke, Balick, & Frank 1992; Mellema 1995; Dwarkadas, Chevalier, & Blondin 1996). In addition, many PNs and pPNs (particularly, but not limited to, those with jets) show point-symmetric morphologies in which nebular features are reflected about the central star (Sahai & Trauger 1998). Such features are difficult to explain with purely hydrodynamic models. More recently, a new problem has emerged as it has been recognized that radiation-wind driving alone can not be responsible for the momentum and energy budgets associated with pPNs (Bujarrabal et al. 2001).

The failure of pure hydrodynamic models has led researchers to consider MHD scenarios for PNs formation. Of particular interest are models in which the magneto-centrifugal forces from an accretion disk (Konigl & Pudritz 2000) both launch the wind and collimate it into a jet. Blackman, Frank, & Welch (2001) and Frank & Blackman (2003) have shown that MHD disk wind models can be quite effective at driving outflows with the total momentum and energy observed in pPNs and PNs. We note that weak magnetic field models are, by their nature, unable to produce observed pPNs momenta, and we do not consider them here (García-Segura et al. 1999).

Of course, any model that requires an accretion disk will, in the context of PNs, require a binary system. Observationally, there exist strong links that argue that disks winds and jets from binary companions play a significant role in pPNs and PNs shaping. Close binary PNs central stars are known to exist (Bond 2000), indicating that interactions between stellar mass-loss processes are likely to occur at some point. More telling is the example of the symbiotic stars. These are binary stars in which a compact companion is known to orbit an AGB star. There are numerous examples of highly collimated outflows from these systems that appear quite similar to PNs (Corradi et al. 2000).

Morris (1987) and Soker & Livio (1994) were the first to identify accretion disks from binaries as the source of highly collimated PNs. A number of studies have explored the formation and properties of these disks (Mastrodemos & Morris 1998, 1999) and the nature of the binary systems that would form them. More recently, Soker & Rappaport (2000) have considered the existence of collimated fast winds (CFWs) from an orbiting companion’s disk and sketched out the flow pattern that would occur as the CFW interacts with the AGB wind. Livio & Soker (2001) applied such a model to M2-9, demonstrating that the resulting flow pattern could, in principle, be understood as the interaction of a CFW and an AGB wind.

In this paper, we use three-dimensional adaptive mesh hydrodynamical simulations to explore the hydrodynamics of AGB wind/CFW interactions. In particular, we focus on explicit realizations of the flow patterns near the two stars in an attempt to understand the limits of this class of models.
Carrying forward simulations such as these pose a number of numerical challenges, and here we attempt to explore only the fundamentals of the hydrodynamics. We note that our models represent the first numerical exploration of this important class of models and should serve to help articulate key issues related to their application and further study.

The structure of our paper is as follows. In § 2, we describe the numerical model, assumptions, and simplifications used in our calculations. In § 3, we describe our results, and in § 4, we present a discussion of their implications as well as our conclusions.

2. NUMERICAL SIMULATIONS

2.1. Numerical Methods and Equations

We work in the regime of ideal gas dynamics in the presence of gravitational sources. Thus we solve the Euler equations, which in conservative form are

\[ \frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0, \]

\[ \frac{\partial \rho \mathbf{v}}{\partial t} + \nabla \cdot (\rho \mathbf{v} \mathbf{v} + p \mathbf{1}) = -\rho \nabla \Phi, \]

\[ \frac{\partial E}{\partial t} + \nabla \cdot (v[E + p]) = -\rho \mathbf{v} \cdot \nabla \Phi, \]

where \( \rho, \mathbf{v} \) and \( p \) are the density, velocity and pressure at a point of the fluid, respectively, \( \mathbf{v} \mathbf{1} \) is the tensor product, and \( \mathbf{1} \) is the unit tensor. Here, \( E \) is the total energy per unit of volume, given by

\[ E = \frac{\rho}{\gamma - 1} + \frac{1}{2} \rho |\mathbf{v}|^2. \]

In addition, we included an advection equation for a passive scalar, which will be used for tracing the different flows. Here, gravity is due solely to the stellar sources, and self-gravity is not included. The gravitational source term is given by

\[ \nabla \Phi = G \left( \frac{M_{\text{AGB}} r}{|\mathbf{r}|^3} + \frac{M_c (r - \mathbf{r}_c)}{|\mathbf{r} - \mathbf{r}_c|^3} \right), \]

where \( M_{\text{AGB}} \) and \( M_c \) are the masses of the AGB and its companion, respectively, and \( \mathbf{r}_c \) is the position of the companion. The gravitational source term is handled in an operator-split fashion via a first-order integration carried out between hydrodynamic time steps.

The numerical simulations are carried out using a version of the YGUAZU-A adaptive code described by Raga, Navarro-González, & Villagrán-Muñiz (2000). This code is well tested and has been used for a variety of supersonic flow problems. The hydrodynamic equations are numerically integrated on an adaptive mesh refined using a second-order flux vector splitting scheme with a Van Leer (1982) algorithm. We used four levels with \( 640 \times 320 \times 320 \) cells on the maximum level. The boundary conditions are transmission for both \( x \) and \( y \) boundaries as well as the maximum \( z \) boundary. A reflecting boundary condition is used on the lower \( z \)-axis, because of the symmetry of the problem. The version of the code used in this work does not include either ionization dynamics or radiative transfer.

We have carried out our simulations assuming an approximately isothermal flow. Thus we choose \( \gamma = 1.01 \). Since our simulations are carried out close to the stars, the densities in the flow are quite high (\( \rho \approx 10^{-18} \text{ g cm}^{-3} \)), and the cooling times will be shorter than what are usually considered in nebular problems. The use of \( \gamma = 1.01 \) produces high compressions and minimal temperature variations behind shocks and is a proxy for the use of a full treatment of time-dependent cooling. In spite of the high densities, the temperature structure behind high-velocity shocks, \( V_s \approx 1000 \text{ km s}^{-1} \), is likely to be complex, and postshock gas may not immediately cool. In such cases, a time-dependent treatment of the cooling should be used. In this paper, we only explore cases with lower velocities for which the cooling timescales will be relatively short compared with dynamical times and the isothermal assumption is acceptable.

2.2. Initial Conditions

Our simulations were designed to explore the interaction of a collimated wind from an orbiting companion interacting with the spherical outflow from the AGB primary. We do not attempt to explore the formation of accretion disks and simply apply an inflow condition at the location at the instantaneous position of the companion. As we discuss below, however, we have chosen binary parameters that are expected to lead to the formation of an accretion disk. We note that these simulations are computationally intensive. We made a number of assumptions and simplifications, which allowed a trade-off between remaining in the correct parameter regime and resolution/computation resource issues.

For the AGB wind, we assume typical values of parameters. We begin with a 1.4 \( M_\odot \) star driving a spherically symmetric mass-loss of rate \( M_{\text{AGB}} = 10^{-6} M_\odot \text{ yr}^{-1} \). We take the radius of the AGB to be a “wind release” radius of \( r_{\text{AGB}} = 4 \text{ AU} \). This is the point at which we inject the wind into the grid with \( v_{\text{AGB}}(r_{\text{AGB}}) > v_{\text{esc}}(r_{\text{AGB}}) \). Given the computational requirements needed for the hydrodynamics alone, we do not include radiation driving on dust in the wind, which would produce a gradual acceleration until a terminal velocity is reached at \( r \gg r_{\text{AGB}} \). Instead, the wind follows a type III solution to the classic Parker wind equation (Raichoudhuri 1998)

\[ (v^2 - c^2) \frac{dv}{dr} = \left( \frac{2c^2}{r} - \frac{GM}{r^2} \right) v, \]

where \( c \) is the sound velocity. Given our boundary conditions at \( r_{\text{AGB}} \), our wind first decelerates because of gravity and then reaccelerates on a trajectory similar to the radiation-driven wind. Thus, the radial structure of the wind close to the star departs from dust-driven wind models. Note that the wind is always supersonic. We have carried out simulations in which we move the inflow boundary further outward such that the wind velocity remains relatively constant and have found no change in the global flow patterns.

For our companion, we chose a 0.6 \( M_\odot \) star with orbital separation \( a = 10 \text{ AU} \). We assumed a circular orbit with a resulting orbital period of \( T_o \approx 22.4 \text{ yr} \). We note that our simulations are carried out with the AGB primary fixed to the origin of the coordinate system. We do not attempt to treat the reflex motions of the companion. This greatly reduces the complexity of the boundary conditions on both the AGB star...
and the jet. Comparison of the various velocities in the problem demonstrates that such an assumption should not change the global flow patterns we attempt to explore. If \( v_1 = (GM_1 \rho)^{1/2} \) is the velocity of the primary about the center of mass, we require a hierarchy of velocities \( v_j \gg v_{AGB} > v_1 \). For our parameters, we find that \( v_j/v_{AGB} = 0.16 \). Thus Coriolis forces will not have a substantial effect on the global flow. In addition, the centrifugal force is at most 10% of the gravity force on the wind. For the present study, we have neglected these terms in the momentum equation, and we expect that “reflex” motions should not have a dramatic effect on our results. Future studies will need a more sophisticated treatment of the boundary conditions and noninertial terms.

Based on the description of Soker & Rappaport (2000), we have chosen to explore two cases for the properties of the jet. For the “weak jet” case, the momentum flux from the AGB star is larger than that from the jet. We define a parameter \( \chi \)

\[
\chi = \frac{M_{AGB} v_{AGB}}{M_j v_j}
\]

where the subscript “\( j \)” refers to properties of the jet. The weak jet models have \( \chi > 1 \), whereas the “strong jet” case has \( \chi < 1 \).

We note that, in general, \( \chi \) depends on the opening angle. However, in this work, the opening angle is fixed for all our models. Since the dynamics of narrow jets interacting with a spherical wind will be more complicated than can be captured by ram pressure or momentum balance arguments, we expect that \( \chi \) will give us only a crude measure of the flow regimes. In the two weak jet simulations we present, we used \( v_j = 200 \text{ km s}^{-1}, M_j = 10^{-7} M_\odot \text{ yr}^{-1} \), and \( M_j = 2.5 \times 10^{-8} \text{ M}_\odot \text{ yr}^{-1} \), which yield \( \chi = 1.25 \) and 5. In the strong jet case, \( v_j = 400 \text{ km s}^{-1} \) and \( M_j = 10^{-7} \text{ M}_\odot \text{ yr}^{-1} \) yields \( \chi = 0.625 \).

We note that values for the parameters taken for the jet are appropriate to an outflow driven by a main-sequence companion for which \( v_j \approx v_{esc} \). We took a main-sequence star as the accretor based on the analysis in Soker & Rappaport (2000). Another possibility is a white dwarf as accretor; however, in this case, the jet velocity is larger, implying a larger cooling time. The isothermal equation of state would not be valid in such a case. In our simulations, the jet is not fully collimated. At the instantaneous position of the companion, we inject a flow for which the angle of collimation is 10°. Note also that our strong jet parameters imply that the companion accretes a large fraction of the AGB wind. This may not be realistic but was chosen to allow us to use a lower jet speed and maintain the isothermal assumption for the computations.

To produce a collimated jet, it is necessary for an accretion disk to form around the secondary star. As discussed in Soker & Rappaport (2000) and elsewhere, a disk can be formed when the specific angular momentum of accreted material \( J_{\text{acc}} \) is larger than the specific angular momentum of a particle in a Keplerian orbit at the equator of the accreting star, \( J_{\text{Ke}} \), i.e., \( J_{\text{acc}}/J_{\text{Ke}} > 1 \). From Soker & Rappaport (2000),

\[
\frac{J_{\text{acc}}}{J_{\text{Ke}}} = 15 \left( \frac{\eta}{0.2} \right) \left( \frac{M_{AGB} + M_\odot}{1.2 M_\odot} \right)^{1/2} \left( \frac{M_2}{0.6 M_\odot} \right)^{3/2} \left( \frac{R_2}{0.01 R_\odot} \right)^{-1/2} \left( \frac{a}{10 \text{ AU}} \right)^{3/2} \left( \frac{v_j}{15 \text{ km s}^{-1}} \right)^{-4},
\]

where \( \eta \) is the ratio of the specific angular momentum of accreted material to that of material that enters the Bondi-Hoyle cylinder (Livio et al. 1986; Soker & Rappaport 2000).

\[ v_j^2 = (v_{AGB}^2 + v_j^2) \]

is the relative velocity of the wind and companion. We take \( \eta = 0.2 \), which applies to a case for which the accretion lies between a fully isothermal and an adiabatic flow and \( R_2 = 0.6 R_\odot \). At \( r = a \), we find a wind speed of \( v_j \approx 15 \text{ km s}^{-1} \). Thus, we find \( J_{\text{acc}}/J_{\text{Ke}} \sim 2.5 \), implying that a disk would form via accretion of the AGB wind. Detailed SPH simulations carried out by Mastrodemos & Morris (1998) using similar parameters as ours tracked the accretion flow and demonstrated that the formation of accretion disks was a robust result. An upper limit to the size of the disk can be calculated using the Bondi-Hoyle radius \( R_d = 2GM_j/v_j^2 = 7 \times 10^{13} \text{ cm} \). A more exact relation for the disk radius is given in Livio & Soker (2001), from which we find \( R_d \approx 1.278 \times 10^{12} \text{ cm} \). This disk size is less than 1 pixel in our simulation (\( \Delta x \approx 0.5 \text{ AU} \) at the highest resolution).

3. RESULTS

In this section, we explore the flow pattern obtained in the weak and strong jet cases and attempt to link these with previous theoretical and observational work.

3.1. Strong Jet

We begin with consideration of the strong jet case whose flow pattern is simpler to visualize and understand. In Figure 1, we present density gray-scale maps of the strong jet simulation in the: Figure 1a, \( xz \)-plane; Figure 1b, \( yz \)-plane; and Figure 1c, \( xy \) plane. The first two cross sections run through the central axis of the computational space, and the third is taken at the base of the flow and includes the AGB and companion/jet inflow boundaries. Figure 1 shows this simulation after 224 yr, or 10 orbital periods.

With a momentum flux ratio of \( \chi = 0.625 \), one would expect that the jet material would be able to propagate through the AGB wind relatively unimpeded. If shocks do occur, they will be found along the sides of the expanding jet column where the jet has pushed AGB material away, entraining and accelerating it. Most importantly, we would not expect the jet to be strongly deflected. The principal modification to the jet’s motion could arise from the shift in its launch point due to the orbital motion of the companion. This will only occur when the orbital period \( T_o \) is comparable to the jet crossing time through the computational domain \( T_{\text{cross}} = H/v_j \) where \( H \) is the height of the computational space. When \( T_{\text{cross}}/T_o > 1 \), then a corkscrew morphological pattern may be expected for the jet’s three-dimensional structure. The spherical AGB wind, on the other hand, will be strongly modified by the passage of jet material. AGB material whose trajectories cross the jet will have to be shocked and diverted. Since we do not expect strong shocks in the jet beam for our strong jet case, the velocity \( v_j = 400 \text{ km s}^{-1} \) is expected all the way up the jet. Thus \( T_{\text{cross}}/T_o = 0.1 \), and we do not expect much “twisting” in the jet morphology.

These expectations are borne out by the simulations. Consider the \( xz \) and \( yz \) cross section in Figures 1a and 1b. Figure 1a shows a broad V-shaped outflow “lobe” driven by the jet as it pushes through the AGB wind. The outer boundaries of the V are shock waves propagating through the expanding AGB wind. This shock creates a dense shell of compressed AGB material (see Fig. 1a). Note that this
structure is not symmetric. The concave shell appears as a darker “bar” on the right side on Figure 1a) with a weaker feature on the left. Note that a similar pattern appears in Figure 1b, but here the jet only appears at a higher value of \( z \), and there are stronger asymmetries between the left and right part of the grid. The asymmetry is a result of the three-dimensional nature of the flow. Shocks are driven through the AGB wind as the jet is swept around with the orbit of the star that produces it. Thus, a given point in the computational space experiences the passage of a succession of strong pressure waves at intervals of approximately \( T_p \). Note also the high-density structure, which can be traced back to the jet source with a second low-density region external to it (the lower left panel of Fig. 1a). This feature may be due to the movement of the jet inflow boundary condition as new cells become injection points while others return to being reflection boundaries.

Fig. 1.—Strong jet gray-scale density maps of cross sections in (a) the \( xz \)-plane, (b) the \( yz \)-plane, and (c) the \( xy \)-plane (i.e., the equatorial plane). Panels on both sides of (c) are enlargements of (a) and (b). The white circles are the position of the center of the stars. The units on the axes are given in AU. The gray scale is the logarithm of the density.
Another nested, V-shaped feature comprised of lighter gray scales is also apparent. This defines the limits of the cavity carved out by the jet. Consideration of the details of the flow patterns show that propagation of the jet through the AGB material results in considerable entrainment of AGB gas (the darker gray scales immediately surrounding the jet) and the final state is one in which the accelerated AGB material fills a large volume of the cone defined by the global flow. This entrainment appears to occur via the oblique shocks that define the interface of the jet/AGB flow. As the jet material is slowed and redirected via these shocks, AGB material is accelerated and mixed into the shell that defines the cavity. Note the periodic structures at the edges of the jet. We conjecture that these features represent Kelvin-Helmholtz (KH) modes, which are only just resolved in these simulations. These modes will enhance AGB acceleration and mixing at the interface of the jet/AGB cavity.

Examination of Figure 1 shows how the orbital motion of the jet drives a positional interchange between the current location of the companion and the location where the strongest jet/AGB wind interaction will occur. In the enlargement of the lower region of Figure 1a, the xy cross section, we see the jet lies to the right of the AGB star. The lower density (lighter gray scale) material can be seen flowing in the z-direction in panel (a). Jet material that has been ejected earlier (and is at higher z) will be farther up the beam and deeper into the three-dimensional volume than the position of this xy cut. The dark gray scales that bound the jet are “sideways” shocks that result whenever a jet flow changes its directions. Such features have been seen and described in simulations of precessing jets (Raga, Cantó, & Biro 1993; Cliffe, Frank, & Jones 1996). The shift in position of the jet and shock features is also apparent in Figure 1b which shows the yz-plane. Here the lower density features that appear above the AGB star are jet material. This structure occurs both because of the opening angle of the jet and its orbital movement. In this simulation, the jet is strong enough that its core remains undisturbed until it leaves the grid. We do see some mild orbital effects in the jet morphology reflected in the apparent change in direction of the jet beam in Figure 1a at the top of the grid where the lightest contours terminate. In the top of Figure 1b, however, the lightest contours represent the undisturbed jet extending to the edge of the grid.

The global morphology of the jet/AGB wind interaction can be seen in Figure 2. This figure shows a three-dimensional volume rendering of the velocity along with an isosurface represented as a wire-frame mesh. Given the limitations of displaying three-dimensional data sets, the reader is encouraged to compare this figure with Figure 1 for a greater comparison of distinct density and velocity figure. Focusing on the velocity, we can clearly see that the jet has only a small trace of a corkscrew pattern and that it is the expansion of the jet and the entrainment of AGB material that contributes most strongly to the velocity (and hence density) pattern. The entrainment of the shocked AGB wind with shocked jet material is particularly apparent in this figure, as the jet is surrounded by lower velocity.
gas. We see shocked/mixed AGB material at a wide range of velocities, with the highest speeds reached being \( v \sim 250 \text{ km s}^{-1} \). Note also the arc of material defined by the wire frame. This feature defines the wake of the jet and companion through the AGB gas via spiral shocks, although it is not possible to distinguish its formation due to gravitational focusing.

The behavior of the flow near the orbital plane displays a number of interesting features. Figure 1c, which shows the base of the computational plane, reveals spiral shocks, which occur because of the gravitational effect of the companion. While our simulation cannot resolve the Bondi radius of the companion and so cannot trap wind accretion, we note that spiral features such as these were also seen in SPH simulations of disk formation (Mastrodemos & Morris 1998). Note that we see that these spiral shocks extend above the equatorial plane in Figure 2, just as Mastrodemos & Morris (1998) did. In our case, however, the spiral shocks above the plane appear closely connected with the wake of the jet. We note that a ring or torus of higher density compressed AGB material appears near the equator. This material is formed via the action of the rotating jet (apparent in Figs. 1 and 2 and predicted by Soker & Rappaport 2000).

In Figure 3, we show one-dimensional cuts of density and velocity through the computational space in the \( xz \) plane at \( x = 80 \text{ AU} \), which corresponds to the center of the AGB star, and \( x = 90 \text{ AU} \), which corresponds to the initial symmetry axis of the jet. The \( x = 80 \text{ AU} \) cut shows a steep rise in velocity, which begins at \( z = 20 \text{ AU} \) and terminates at \( z = 80 \text{ AU} \). This feature reflects the acceleration of the AGB material as it interacts with the jet. For \( z > 80 \text{ AU} \), the line cuts through jet material traveling close to its injection velocity. The \( x = 90 \text{ cut} \) shows a velocity that is approximately constant, reflecting the fact that the jet is strong enough that its core remains undisturbed by the AGB material. Note the oscillations in density and velocity that reflect the apparent KH modes along the jet/AGB interface.

To summarize, the interaction of the jet with the slow wind produces a complex morphology composed of multiple shock waves driven into both the AGB wind and the jet. In spite of its complexity, a quasi-steady state pattern does emerge in which the entire flow is embedded in a V-shaped shock propagating outward through the AGB wind. The shock produces a dense shell of compressed AGB material. Within this shell is a cavity produced by the propagating jet, which does not assume a strong corkscrew flow pattern because the jet transit time is shorter than the orbital period, i.e., jet material moves fast enough that it leaves the grid in a time less than an orbital period, \( t < T_o \). Significant entrainment of AGB and jet material occurs because of shocks at the jet/AGB interface. At the equator, spiral shocks are driven into the AGB wind because of the orbital motion of the companion and the effect of the jet.

### 3.2. Weak Jet

We now consider the simulation results for the two cases in which momentum in the jet is weaker than that in the

![Fig. 3.—Cuts in velocity and density. The left top panel is the velocity along a vertical line at \( x = 80 \text{ AU} \) (the center of the AGB star). The bottom left panel is the density along the same line. The right panels are taken at \( x = 90 \text{ AU} \) (approximately the position of the jet).](image-url)
AGB wind. As the AGB wind momentum flux becomes successively larger than that in the jet, we expect strong shocks to form in the jet beam, as well as jet deflections caused by the ram pressure of the AGB material. The first weak jet simulation we present has \( \chi = 1.25 \), and density maps in different orthogonal planes are shown in Figure 4. The second case has \( \chi = 5 \) and is shown in Figure 6. Once again, both simulations correspond to an age of 224 yr, or 10 orbital periods.

When the flux of momentum of the AGB wind is slightly higher than that in the jet, \( \chi = 1.25 \), the V-shaped shocks in the AGB wind marking the boundary of the interaction region still appear. However, as shown in Figure 4, the effects of the stronger AGB wind are apparent. First, it is possible to see a shallow inclination of the jet in Figure 4a (see left panel detail). For this value of \( \chi \), the deflection of the entire jet does not occur. Instead, we see a narrowing of the jet on the side facing the AGB in the region at low \( z \).

![Fig. 4.—Weak jet gray-scale density maps of cross sections in (a) the \( xz \)-plane, (b) the \( yz \)-plane, and (c) the \( xy \)-plane (i.e., the equatorial plane) with \( \chi = 1.25 \). Panels on both sides of (c) are enlargements of (a) and (b). The gray scale is the logarithm of the density.](image-url)
close to the jet source. In the language of Soker & Rappaport (2000), the deflection angle of the jet is still smaller than the collimation angle for these parameters. Thus, only the collimation on one side of the jet is affected by the AGB ram pressure. Thus, we see that our expression for $\chi \approx 1$ does give some measure of the boundary between strong and weak jet regimes even after accounting for the opening angle of the jet.

What is more apparent in this case, as compared to the strong jet simulation, is the strong shock in the jet beam. We see a shock and dense shell in Figure 4a) at $z = 100$ AU. Note that the apparent termination of the jet at $z = 100$ AU in Figure 4a and $z = 280$ AU in Figure 4b are evidence of its three-dimensional structure. In Figure 4a, for example, jet material that had been ejected earlier (and is now higher at $z = 100$ AU) will be further up the corkscrew and deeper into the three-dimensional volume than the position of this $xy$ cut. The dynamically important connection is that the shock decreases velocities in the material surrounding the jet. The $T_{\text{cross}}/T_0$ becomes larger, and the orbital motion begins to exert a greater influence on the jet morphology in terms of creating a global “corkscrew” pattern.

Near the equatorial plane and close to the AGB star, the flow structure is very similar to the strong jet case. Above the AGB star, there is a shock wave and a thin shell of compressed material. This shock can be better appreciated in the one-dimensional plots in Figure 5. Here the shock is clearly seen at $z \approx 20$ AU at the $z$-axis ($x = 80$ AU) and in the $yz$-plane (left panels in Fig. 5). After this shell, the AGB material is gradually accelerated via entrainment reaching a velocity of $\approx 180 \text{ km s}^{-1}$. Using this value in the expression for the crossing time, we find $T_{\text{cross}}/T_0 \approx 0.4$.

When $\chi$ increases, the deflection angle should increase (see Soker & Rappaport 2000, although their expression for the jet bending only applies when the bend occurs on scales smaller than the orbital separation $a$). Thus, for the $\chi = 5$ simulation, we expect greater influence of the AGB ram pressure on the jet beam. This expectation is borne out in the $xz$ and $yz$ maps in Figure 6, which clearly show that the entire beam has bent away from the AGB star. Globally, the cross sections show a succession of staggered “donkey ear” shaped shocks and dense shells propagating through the AGB wind. Such a pattern was predicted by Soker & Rappaport (2000) and occurs because of the deflection (and disruption) of the jet by the slow wind over an entire orbital cycle.

Consideration of the postshock velocities in the jet and surrounding material (Fig. 7) shows that $T_{\text{cross}}/T_0 \approx 0.86$. With this in mind, the enlargement of the $xy$-plane map in Figure 6 allows us to see the origins of the large-scale flow pattern. The jet can be seen emanating from the companion on the right side of the AGB star in Figure 7a. Unlike the strong jet case, however, the jet only propagates to $z \approx 30$ AU before it is shocked via its interaction with the AGB wind. The jet is bounded above by a dense shell comprising both shocked jet and shocked AGB material. The continuation of this structure can be seen in the $yz$ cross section on the left side of the AGB star.

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**Fig. 5.** Cuts in velocity and density for a $\chi = 1.25$ simulation. The left top panel is the velocity along a vertical line at $x = 80$ AU ($z$-axis). The bottom left panel is the density along the same line. The right panels are similar, but they are taken at $x = 90$ AU (approximately the position of the jet).
On the $xz$- and $yz$-planes, it is possible to see three “ear shaped” lobes, two on Figure 7a and one on Figure 7b. Each of these features is part of a continuous three-dimensional structure. What we see in the cuts are cross sections of shocks formed from the interaction of the jet and AGB wind, each initiated at successive $\frac{1}{4} T_o$ time delays. Note only three lobes are apparent in the figure, implying that jet material ejected in the first $\frac{1}{4}$ of the current orbital period has already left the grid. Unlike the strong jet case, some of the jet material exits from the side of the grid because of the deflection.

The three dimensional structure of the flow is shown in Figure 8 for the $\chi = 5$ case. The isosurfaces in density have values $6.168 \times 10^{-18}$, $5.607 \times 10^{-19}$, and $1.78 \times 10^{-19} \text{g cm}^{-3}$, respectively. While this figure is complex and the identity of the isosurfaces can be lost (which is in itself a measure of the nature of the nonlaminar flow in this case), it is helpful to note that lower density isosurfaces are located at higher $z$. This figure
demonstrates the formation of the global corkscrew flow pattern produced by the orbital motion of the source. Note that we see less than a single turn of the corkscrew as would be expected for the value of $T_{\text{cross}}/T_o$. The complexity of the flow is also apparent in this figure. The action of multiple shocks passing through the AGB wind and jet creates a mix of structured and disordered isodensity surfaces. A "donkey-ear" shock is apparent in the middle and lower regions of the computational space. Note also that the spiral shocks in the AGB wind and the global $V$-shaped (in cross section) shock that bounds the jet/AGB wind interaction are also apparent in this figure. We note that explorations of figures such as these show that the exact choice of isodensity surfaces does not change the qualitative conclusions that can be drawn from this figure.

4. DISCUSSION AND CONCLUSIONS

We have presented new simulations of the interaction of a slow AGB wind with a collimated fast wind (CFW) driven by a binary companion. Our binary parameters were selected such that an accretion disk would form even though we are unable to resolve such a flow pattern. We report three simulations segregated by the ratio of the AGB wind momentum flux to the CFW momentum flux (denoted $\chi$). We simulated a strong jet case, $\chi = 0.625$, and two weak jet cases, $\chi = 1.25$ and 5.0. The binary period was $T_o = 22.4$ yr, and we carried out all simulations for at least $10T_o$.

We find that $\chi$ does provide a measure of prediction for differentiating the behavior of the simulations, although future studies will be needed to account for the effect of the opening angle of the jet. Strong jets are able to propagate off the grid ($Z_{\text{max}} = 320$ AU) without severe deflections or disruptions. When $T_{\text{cross}}/T_o \ll 1$ as it was in the strong jet case, the orbital motion does not affect the global morphology of the jet. When the jet becomes progressively weaker and $\chi$ becomes larger, we see a trend toward both stronger deflections (the jet is bent away from the direction of the AGB star) and stronger disruption (strong shocks bounding the jet beam). The location in height of the first shock moves closer to the jet source as $\chi$ increases. Weaker jets also lead to more complicated global flow patterns, with features such as multiple "donkey ear" shaped lobes appearing at well-characterized intervals in height.

Our simulations support the idea that collimated jets formed close to the central source are the agents shaping some PNs. This paradigm has steadily been gaining favor. Morris (1987) and Soker & Livio (1994) both proposed that jets could be formed via accretion disks in PNs. Sahai & Trauger (1998) proposed that jets formed very early in the PNs or pPNs phase were responsible for the bulk of the nebular shaping with the fast wind from the central star of a PNs merely burnishing the details of the nebular shapes. More recently, Soker (2002) and Lee & Sahai (2003) have both explored hydrodynamic models of jets driving through circumstellar environments. Our work is complementary to Lee & Sahai (2003) in that we examine the flow pattern on smaller scales. Soker & Rappaport (2000) investigated both the hydrodynamics and population statistics of the CFWs and the outflows they would drive in binary systems. They concluded that narrow-waisted PNs are likely the result of CFW/AGB wind interactions. Livio & Soker (2001) applied such a model to M2-9, concluding that a weak
jet with a severely deflected wind was responsible for both the outflow shape and the shadowing of ionizing radiation from the hot companion.

Our results are most directly relevant to the work of Soker & Rappaport (2000) and can be seen as a test of the ideas put forth there. We find that much of the flow pattern predicted in that work is obtained in our simulations. We do see some evidence for the creation of a region of enhanced density along the equator due to the action of the CFW as predicted by Soker & Rappaport (2000). The successive donkey-ear lobes seen in the simulations were anticipated by Soker & Rappaport (2000) as well. It is not yet clear if their strong deflection flow pattern will be obtained in real systems, however, since we are currently not able to resolve details of the jet flow when shocks form very close to the inflow source. We will take up this issue, along with the effects of wind acceleration, in a future study.

Our most important conclusion, however, is that CFWs do appear as good candidates for creating narrow-waisted nebula. The strong jet case, in particular, shows that the effect of a CFW is primarily confined within a V- or U-shaped shock in the expanding AGB wind. Such a jet will drive a very narrow-waisted flow as there is no expansion along the equator. Thus, on large scales, the lobes driven by the jet will appear to pinch severely at the outflow source. We note that recent Hubble Space Telescope (HST) observations of the near nuclear regions of M2-9 appear to confirm the flow patterns we see in the strong jet case (B. Balick 2003, private communication) with oppositely oriented V-shaped features extending from the unresolved central source out to scales of $10^{15}$ cm.

This work constitutes a first attempt at mapping out the flow dynamics of outflows driven via interactions between an AGB wind and a jet from a binary companion. Here we have attempted to examine the global properties of the outflows on scales ranging from 5 to 300 AU. While the use of an AMR code allowed us to capture details such as shocks in the AGB wind and jet beam, the resolution used and limited physics included still leave many open questions. These include the nature of the flow when $\chi$ is very large, the role of radiation driving of the AGB wind, the effect of ionizing radiation, and the effect of magnetic fields (which should be present if the jets are magneto-centrifugally launched). In spite of these limitations, our simulations provide further support for the argument that jets driven by accretion disks can explain many features observed in bipolar PNs.

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