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DISKS AND JETS IN PLANETARY NEBULAE

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ABSTRACT

Highly collimated bipolar outflows and jets have been recently observed to be associated with some planetary nebulae. Assuming that the mechanism of jet formation requires the presence of an accretion disk, we examine the possibilities to form such disks in the case of planetary nebulae with binary nuclei. We show that at the emergence from the common envelope phase, some systems can form accretion disks.

Planetary nebulae therefore form the newest addition to the classes of objects (which so far include active galactic nuclei, young stellar objects, and some X-ray binaries) which produce jets.

Subject headings: accretion, accretion disks — ISM: jets and outflows — planetary nebulae: general

1. INTRODUCTION

Highly collimated jets are associated with a number of astrophysical objects. The best studied examples are the radio and optical jets observed in active galactic nuclei (e.g., Boksenberg et al. 1992; Biretta, Stern, & Harris 1991) and the optical jets observed in young stellar objects (e.g., Reipurth & Heathcote 1993).

In spite of extensive research and many suggested models, the exact mechanism of jet formation and collimation is still unclear (see, e.g., Blandford 1993; Pringle 1993 for reviews).

Other types of less collimated bipolar flows are associated with many nebulae (e.g., He 2-104, BI Cru, My Cn 18, Corradi & Schwarz 1993; OH 17.7 − 2.0, La Bertre 1986; IRAS 18059 − 3211, Ruiz et al. 1987; NGC 3242, Soker, Zucker, & Balick 1992; R Aqr, Burgarella & Paresce 1991, 1992).

One of the models for the formation of these bipolar flows involves the action of a binary companion to the central star (e.g., Morris 1987, 1990; Soker 1992b; Livio 1992). In particular, it has been shown that the axially symmetric morphology of planetary nebulae (PNs) with close binary nuclei is consistent with the predictions of common envelope (CE) evolution, when these are coupled with the "interacting winds" model (Bond & Livio 1990). The following sequence of events is supposed to occur. During the CE phase, mass is ejected preferentially in the orbital plane (Taam & Bodenheimer 1989; Livio & Soker 1988). This forms a "density contrast" between the equatorial plane and polar directions. When the hot central star is exposed, it emits a fast, dilute wind (Kwok 1982; Kahn 1982). The fast wind catches up with the slowly moving material and shocks it. It has been shown that this interaction can produce the observed morphologies of PNs (e.g., Balick 1987; Soker & Livio 1989; Soker 1990; Icke, Balick, & Frank 1992a; Frank et al. 1993).

A prominent feature in many elliptical PNs is the presence of two opposite bright knots ("ansae") along the symmetry axis of the PN. In some cases a "jet"-like structure seems to connect the ansae to the central star (e.g., NGC 3242, Balick 1987). Similarly, Gieseking, Becker, & Solf (1985) interpreted their observations of NGC 2392 (the "Eskimo" nebula) as a jetlike bipolar outflow extending from the central star to the two ansae. They found the angle of collimation in NGC 2392 to be less than 10°, and the bipolar flow was found to have a velocity of 190 km s^{-1} and a total mass of 2.5×10^{-4} $M_\odot$. Soker (1990, 1992a) proposed that in these PNs, "jets" are formed in the period between the end of the asymptotic giant branch (AGB) wind and the beginning of the fast wind. He regarded the fact that in some cases (e.g., NGC 3242 and NGC 7009) the morphology and velocity map suggest that the "jets" are bent, as providing support for his proposed scenario.

The above-mentioned works are related mostly to "classical" PNs in which the jets, and the features associated with them, contain a very small fraction of the mass (< 1%) and volume of the nebula. Recently, a few PNs were observed to exhibit well-collimated jets, (in some cases the jet and its "plume" seem to compose the main part of the nebula). These include K1-2 (Bond & Livio 1990; Lutz & Lame 1989), a few of the objects listed by Schwarz, Corradi, & Melnick (1992), which exhibit a point symmetry (e.g., IC 4634), and possibly the "Calabash Nebula" (Icke & Preston 1989). This finding, together with the less collimated "jets" in the "classical" PNs, provides us with the exciting possibility of studying the mechanism of jet formation and collimation in a new type of object.

In the present, still preliminary work, we examine three possible scenarios which can (in principle at least) lead to the formation of PN jets. More specifically, we examine the possibility to form accretion disks in these systems. Two of these scenarios have been mentioned previously but have never been studied in detail. We quantify somewhat both scenarios. In § 2 we discuss the possibility to form jets in binaries in a pre–CE phase. In § 3 we calculate some of the properties of the accretion phase (of backfalling material) which can follow the fast wind phase. We find that this process may be more relevant to the evolution of hot white dwarfs. In § 4 we discuss the post–CE phase. A general discussion and summary follow (§ 5).

2. PRE–PLANETARY NEBULA JET?: SECONDARY OUTSIDE THE ENVELOPE

One of the possibilities to form jets (in principle at least) involves a configuration in which a binary companion revolves around an AGB star. Morris (1987) suggested that a secondary

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1994ApJ...421..219S

star, which orbits outside the primary AGB star, can accrete from the AGB star and form jets through the formation of an accretion disk. Since he, however, did not study this process in detail, we will attempt to quantify this process somewhat in this section.

Because of the deep convective envelope of an evolved AGB star, if the primary is more massive than the secondary ($M_1 > M_2$) at the onset of Roche lobe overflow (RLOF), the system will enter a CE phase. Taking into account more exactly the response of the star to mass loss, Pastetter & Ritter (1989) find that the constraint at the onset of RLOF is even stronger, a CE resulting if $M_1 > 0.6\ M_2$. However, the primary, which is initially the more massive star, can become the less massive by extensive mass loss. Most of the mass loss occurs during the earlier AGB phases. In order for the primary to evolve to an AGB star and reach a situation in which $M_1 < 0.6 M_2$ at the onset of RLOF, Pastetter & Ritter find that the initial mass ratio $q_i \equiv (M_1/M_2)_i$ must be in the range $1.2 \lesssim q_i \lesssim 1.8$.

The allowed range for q_i has some interesting consequences for these systems. Since the cores of PNs have a mass of typically $\sim 0.6\ M_\odot$, and since we require $M_1 < 0.6 M_2$ at the onset of RLOF, the secondary must have a mass satisfying $M_2 > 1\ M_\odot$, and the initial mass of the primary must be $M_{1i} > 1.2\ M_\odot$. In order to be left with a substantial envelope mass on the primary at the onset of RLOF (and assuming a nebular mass of order $0.2\ M_\odot$), the limit is even higher $M_{1i} > 1.4\ M_\odot$. More typically, initial primary masses will probably be $M_1 \gtrsim 2\ M_\odot$. This implies that these systems will be more concentrated toward the galactic plane than PNs in general. This is compatible with the observational finding of Zuckerman & Gatley (1988) that PNs with higher degree of asymmetry tend to be closer to the galactic plane (assuming that such systems will indeed form nonspherical nebulae).

For mass transfer to take place, the distance between the centers of the two stars has to be $\lesssim 10^3\ R_\odot$, and with the masses quoted above, the primary orbital velocity and the orbital period are $v_1 \gtrsim 10$ km s^{-1} and $P \sim 10$ yr, respectively. Prior to RLOF, the AGB star will lose mass in a wind, with a typical velocity of $v_w \sim 10$ km s^{-1}. The near equality between the orbital velocity and wind velocity can lead to a complicated mass flow. In the equatorial plane, for example, we expect shocks to develop, when material blown in the direction of the primary motion collides with material blown half a period earlier, at a velocity of $v_w - v_1$. As the primary continues to evolve along the AGB, it approaches its Roche lobe. Even before filling the lobe, the wind will start to be collimated toward the secondary (Friend & Castor 1982). This can result in an accretion disk around the secondary, even before RLOF.

Let us examine the total mass of the nebula that could result from such a system. At the onset of RLOF we have $M_1 < 0.6 M_2$. Since the core of the primary AGB star has a mass of greater than $0.6\ M_\odot$, the total mass available after the onset of RLOF (and therefore for the jet) is

$$\frac{M_{\rm jet}}{M_\odot} < 0.6(M_2/M_\odot - 1)\ . \quad (1)$$

If before the AGB phase the star loses $\sim 0.2\ M_\odot$ (Harpaz & Shaviv 1992), the total mass which is blown during the AGB phase is

$$M_{\rm AGB} \simeq q_i M_2 - 0.8\ M_\odot\ . \quad (2)$$

From equations (1) and (2) we see that there is in fact more mass in the nebulae formed prior to the onset of RLOF than in the one formed after the RLOF. In practice, however, the nebulae of PNs contain less than $\sim 1\ M_\odot$.

There is nothing in the above scenario (and indeed in the original suggestion of Morris 1987) to indicate that *the formation of a jet is inevitable.* In fact, our entire discussion merely gives certain ranges in parameter space for which *disk formation is plausible.* The question of the exact conditions under which disk accretion produces jets is at present still entirely open. In fact, it may be that some of the configurations that are supposed to lead to an unstable mass transfer (see above) produce jets rather than (or before) a CE configuration. The jet observed in SS 433 (in which a compact object accretes at a critical rate, see, e.g., Vermuelen 1993 for a review) can be regarded as providing support for such a scenario.

3. POST–PLANETARY NEBULA JET?: ACCRETION FROM THE NEBULA

We first examine the conditions under which the central star (CS) accrete mass from the nebula. In a spherical geometry, in which the postshock fast wind does not cool, there is a hot bubble around the central star of density

$$\rho = \frac{M_b}{4\pi r_1^3/3} = 1.3 \times 10^{-26}\left(\frac{M_b}{10^{-4}\ M_\odot}\right)\left(\frac{r_1}{0.5\ {\rm pc}}\right)^{-3}\ {\rm g\ cm^{-3}}\ , \quad (3)$$

where M_b is the total mass in the hot bubble and r_1 is the outer radius of the bubble. For a typical velocity of 1000 km s^{-1}, the fast wind postshock temperature is $T_f \sim 10^8$ K, and the radiative cooling time of the bubble will be

$$\tau_{\rm cool} \sim 2 \times 10^{10}\left(\frac{\rho}{10^{-26}\ {\rm g\ cm^{-3}}}\right)^{-1}\left(\frac{T_f}{10^8\ {\rm K}}\right)^{1/2}\ {\rm yr}\ . \quad (4)$$

In order for the matter from the nebula to flow back, the hot bubble must cool through heat conduction, or the hot gas should leak through the nebula. For nonspherical nebulae the latter is more likely, since magnetic fields are likely to suppress heat conduction.

We assume that the fast wind ceases when the gravity of the white dwarf reaches $\log g \simeq 5$ (cm s^{-2}) (a radius of 0.5 $R_\odot$; Volk & Kwok 1985) and that several hundred years later, the pressure of the fast wind on the inner boundary of the nebula becomes negligible. We further assume (for simplicity) that the nebula shell has a constant density ρ_0, that it is thin, and that it expands at a constant velocity v_0. We assume spherical symmetry and mark the inner shell boundary by r_1. With these initial conditions, a rarefaction wave will start propagating into the nebula from r_1. We use the self-similar solution of a rarefaction wave in one dimension (e.g., Landau & Lifshitz 1987). The rarefaction wave moves with the sound speed C_s relative to the shell, and with $v_0 + C_S$ relative to the CS. The inflow velocity of the inner boundary of the nebula relative to the CS is

$$v_b = v_0 - 2C_s/(\gamma - 1) = C_s[\mathcal{M} - 2/(\gamma - 1)]\ , \quad (5)$$

where γ is the adiabatic index and $\mathcal{M}$ is the Mach number of the shell. At a time t the rarefaction wave and the inner boundary of the inflowing matter are at

$$r_a = (C_s + v_0)t + r_i\ ;$$
$$r_b = C_s[\mathcal{M} - 2/(\gamma - 1)]t + r_1\ , \quad (6)$$

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1994ApJ...421..219S

respectively. In the spherical geometry under study here, $r_b < 0$ means that the matter has been accreted (or at least is available for accretion) by the CS. The accretion will occur only if $\mathscr{M} < 2/(\gamma - 1)$ and will start at a time $t_{acc} = r_1/C_s[2/(\gamma - 1) - \mathscr{M}]$. Substituting the velocity $-2C_s/(\gamma - 1)$ in the density equation (eq. [99.15] of Landau & Lifshitz 1987), and then integrating the density over a volume extending from $r_b < 0$ to $r = 0$, gives the total mass available for accretion. We substitute for the density in the shell a value based on $4\pi r_1^2 v_0 \rho_0 = \dot{M}_s$, where $\dot{M}_s$ is the mass-loss rate of the AGB progenitor. The mass available for accretion, as a function of time, is then

$$M_{acc} = \dot{M}_s t \frac{1}{\mathscr{M}} \left(\frac{\gamma - 1}{\gamma + 1}\right)^{(\gamma+1)/(\gamma-1)} \times \left(\frac{2}{\gamma - 1} - \mathscr{M} - \frac{r_1}{C_s t}\right)^{(\gamma+1)/(\gamma-1)}, \quad t > t_{acc} . \quad (7)$$

We can obtain a more transparent expression for the behavior immediately after accretion starts if we write the last equation in terms of the dimensionless time $\tau = (t/t_{acc} - 1)$,

$$M_{acc} = \dot{M}_s t_{acc}(\tau + 1) \frac{1}{\mathscr{M}} \left(\frac{\gamma - 1}{\gamma + 1}\right)^{(\gamma+1)/(\gamma-1)} \times \left[\frac{\tau[(2/\gamma - 1) - \mathscr{M}]}{\tau + 1}\right]^{(\gamma+1)/(\gamma-1)}, \quad \tau > 0 . \quad (8)$$

From expression (8) we see that at the beginning the accretion rate (or the mass available for accretion) increases very fast, faster than τ^4. The accretion rate asymptotically reaches the value of

$$\dot{M}_{acc} = \dot{M}_s \frac{1}{\mathscr{M}} \left(\frac{\gamma - 1}{\gamma + 1}\right)^{(\gamma+1)/(\gamma-1)} \times \left(\frac{2}{\gamma - 1} - \mathscr{M}\right)^{(\gamma+1)/(\gamma-1)}, \quad \tau \gg 1 . \quad (9)$$

For $\gamma = 5/3$, $\mathscr{M} = 2$ we get from equation (9) $\dot{M}_{acc} = \dot{M}_s/512$. With these values and with $C_s = 15$ km s^{-1} and $r_1 = 10^{18}$ cm, we find $t_{acc} = 2 \times 10^4$ yr. Thus the accretion phase will typically start after the main nebula has faded. The accretion rate will be small compared with the mass-loss rate of the AGB star, but can be in the range of $\dot{M}_{acc} \simeq 10^{-7}$ to 10^{-9} $M_\odot$ yr^{-1}. Since the nebula shell is finite in mass, the rarefaction wave will eventually decay. If the accretion process continues for a few t_{acc}, the total mass accreted can be $\lesssim 10^{-2}$ $M_\odot$. However, the exact amount will depend sensitively on the accretion rate, which will determine if the material can burn stably (e.g., Nomoto 1982).

It is possible that at very early times a shock will form around the CS. The shock front will then propagate outward. We find that for shock radii larger than $\sim 10^{13}$ cm the shocked material has enough time to cool as it flows inward, and thus the accretion process can continue. At such small radii the acceleration by the CS must be taken into account.

What is the form of the accretion in a case in which the CS is orbiting a close companion? This depends on the specific angular momentum of the backflowing material, with respect to the primary star. In cases in which the system went through a CE phase, for example, the nebula is likely to have high specific angular momentum. As it flows back, the matter can form a disk around the compact CS. Under some circumstances that are not well understood yet, the disk might form jets. Predicting the specific angular momentum of the backflowing material, and calculating the exact form of the accretion, are beyond the scope of the present paper. If jets are indeed formed, we expect them to have the following properties. They will exist long after the nebula started fading, though they can exist before the nebula completely dissipated into the interstellar medium (ISM). The escape velocity from the central star is $\sim$ several times 1000 km s^{-1}, which will also be a typical jet velocity. The total mass in the jet will be less than 10^{-3} $M_\odot$. With a mass flow rate of 10^{-6} $M_\odot$ yr^{-1}, and an opening angle of 5°, the density in the jet will be more than two orders of magnitude lower than in a typical nebula. The surface brightness will thus be more than four orders of magnitude lower. Such jets will therefore be very difficult to detect.

The scenario described above may result in a morphology that is similar to that of the almost spherical PN A30. A30 has a spherical smooth halo of radius 100″, with an extremely small deviation from sphericity on its edge (Balick 1987). It has two knots at a distance of 14″ from the center, which may result from jets (like ansae in other PNs). The sense of the asymmetry in the knots and the edge of the nebula is about the same. The absence of a rim in this object indicates that the fast wind ceased a long time ago. According to the scenario presented in this section, a slow rotation of the AGB star (of the order of a few percent of the breakup speed), which may have resulted from a low-mass secondary, caused the tiny deviation from sphericity. The specific angular momentum associated with the rotation of the AGB star is sufficiently high to form a disk around the compact central star.

Before ending this section, we would like to mention two points: (i) a backflow followed by an accretion event was suggested by Chevalier (1989) for the case of SN 1987A. (ii) The accretion of even a very small amount of nebular gas can form a white dwarf (WD) with a relatively high hydrogen abundance. This may have important implications for models of WD cooling, which try to explain the hydrogen abundances at different phases on the cooling track (see e.g., Sion 1986).

4. POST–COMMON ENVELOPE JET?: EMERGENCE FROM THE COMMON ENVELOPE

The PN K1-2 exhibits what appears to be a well-collimated jet (Bond & Livio 1990; Lutz & Lame 1989). In this case, the central star, VW Pyx, is a known close binary with an orbital period of 0.676 days (Bond 1988). It is clear that for this system to form the binary had to undergo CE evolution (see, e.g., Livio 1992 for a recent review). Some of the objects which exhibit a point symmetry and are discussed by Corradi & Schwarz (1993) and Schwarz et al. (1992) are probably similar.

The existence of PNs with close binary nuclei, in which the jet and its associated plumes appear sometimes to represent a large fraction of the observed nebula, raises the following important question: how can such a jet form in a post–CE binary?

The "interacting winds" model, coupled with predictions of CE evolution, has proved very successful in generating bipolar flows (e.g., Balick, Preston, & Icke 1987; Soker & Livio 1989; Icke et al. 1992a; Frank et al. 1993). In this scenario, the density contrast in the cold, slow wind, between the equatorial and polar directions (which can be generated by CE evolution), serves to give inertial collimation to the hot, dilute wind, which is emitted by the exposed central object. Inertial effects

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1994ApJ...421..219S

("funnels") can be generally expected to lead to moderate collimation. Highly collimated jets have actually been obtained in calculations which included both inertial collimation and the effects of the inner oblique shocks (Icke et al. 1992). However, there still appear to be two main problems for inertial collimation (alone) to form highly collimated jets. (i) It requires a profile of the density contrast which appears to be difficult to achieve. (ii) In order to obtain long-lived jets, it is necessary to continuously provide pressure to the confining material (e.g., Blandford & Rees 1974; Königl 1982). This does not appear easy in the context of PNs with binary nuclei, where the profile of the density contrast can change (due to material motion) on a relatively short time scale.

Another mechanism of collimation involves the formation of de Laval nozzles (if the flow is adiabatic) or the formation of an elongated cavity (in the highly nonadiabatic case). Such mechanisms have been suggested to operate in the case of young stellar objects (e.g., Raga & Canto 1989; Königl 1982). More recently, however, serious objections to this mechanism have been raised (in particular because of the small scale on which the collimation occurs, see e.g., Blandford 1993 for a review).

Consequently, in spite of the fact that the exact mechanism by which jets are produced in AGNs and young stellar objects is not entirely clear yet, there seems to be a general consensus about the necessity to have two essential ingredients: (i) an accretion disk, and (ii) a vertical magnetic field threading the disk. The first ingredient provides the axial symmetry and the energy source, and the second is supposed to play a role in the acceleration and collimation (e.g., Blandford 1993; Pringle 1993).

Assuming that at least the most basic ingredients of the mechanism of jet production are more or less universal, we must therefore look for the possibility to form an accretion disk in the final stages of the CE phase. To this goal, it is instructive to follow the evolution of the secondary star during the spiraling-in process. We will use the numerical results obtained by Hjellming & Taam (1991).

In the initial spiraling-in phases, the secondary accretes nonadiabatically. Because of the relatively high specific entropy of the CE material (compared to that at the secondary's surface), the secondary expands as a result of accretion. Expansions by a factor of up to 90 were obtained by Hjellming & Taam (1991) (depending on the mass of the AGB star and its evolutionary state, see also Prialnik & Livio 1985). As the spiraling-in process continues, the radius of the secondary at some point starts to exceed the radius of its tidal lobe. From this point on, the secondary can start losing mass, back to the CE. This process was modeled by Hjellming & Taam approximately, by imposing the condition $R_{\rm secondary} = R_{\rm RL}$, where $R_{\rm RL}$ was the radius of the secondary's Roche lobe. High mass-loss rates were thus obtained for a relatively short, nearly adiabatic phase. In the sequences calculated by Hjellming & Taam, which involved a 1.25 $M_\odot$ secondary, most of the previously accreted mass was returned to the CE, with the secondary retaining typically less than 0.01 $M_\odot$.

The most significant outcome of these calculations, for our present purposes, is the fact that *upon emerging from the CE the secondary departs significantly from thermal equilibrium.* In order to demonstrate that this result, which was obtained in the calculations of Hjellming & Taam, is indeed quite general, we will now show that the final stages of the CE phase are of extremely short duration.

The orbital decay is governed by the equation (as long as the local density is lower than the mean density interior to the orbit)

$$\frac{GM_r M_2}{2r^2}\frac{dr}{dt} \simeq -L_{\rm drag}\,, \tag{10}$$

where M_r is the mass interior to radius r in the AGB-type configuration, M_2 is the secondary mass, and $L_{\rm drag}$ is the rate of energy dissipation by the frictional drag, which is given approximately by (e.g., Shima et al. 1985)

$$L_{\rm drag} \simeq \pi R_A^2 \rho v_{\rm orb}^3\,, \tag{11}$$

where R_A is the accretion radius $R_A \simeq 2GM_2/v_{\rm orb}^2$, ρ is the local density in the CE, and $v_{\rm orb}$ is the orbital velocity. Expression (11) can overestimate the rate of drag energy deposition when R_A becomes of the order of a density scale height or larger. For our present, limited purposes, however, we will take the secondary to be of a very low mass ($M_c/M_2 \sim 10$, where M_c is the core mass of the AGB star), in which case expressions (10) and (11) are adequate.

The density profile in the convective part of AGB stars can be quite well approximated by an expression of the type (e.g., Soker 1992a)

$$\rho(r) \simeq \rho_0\left(\frac{M_{\rm env}}{M^0_{\rm env}}\right)\left(\frac{R_{\rm star}}{R_0}\right)^{-1}\left(\frac{r}{r_0}\right)^{-2}, \tag{12}$$

with typical values being $\rho_0 = 1.5 \times 10^{-8}$ g cm^{-3}, $M^0_{\rm env} = 0.1$ $M_\odot$, $R_0 = 400\,R_\odot$, $r_0 = 100\,R_\odot$.

Substituting relations (11) and (12) into (10) allows for a fully analytic solution for the separation as a function of time of the form (we have taken $t = 0$ to be at $r = r_0$)

$$F(r/r_0) - F(1) = -Ct/\tau_0\,, \tag{13}$$

where

$$F(X) \equiv \left(A + \frac{D + 2BX}{4}\right)[X(D + BX)]^{1/2} + \frac{D}{B^{1/2}}\left(A - \frac{D}{4}\right)\ln\left[B^{1/2}X^{1/2} + (D + BX)^{1/2}\right]\,, \tag{14}$$

and the parameters A, B, C, D, and τ_0 are given by

$$\begin{aligned} A &= \frac{M_c}{M_\odot} - \frac{4\pi r_0^3\rho_0}{M_\odot}\left(\frac{R_c}{r_0}\right)\left(\frac{M_{\rm env}}{M^0_{\rm env}}\right)\left(\frac{R_{\rm star}}{R_0}\right)^{-1},\\ B &= \frac{4\pi r_0^3\rho_0}{M_\odot}\left(\frac{M_{\rm env}}{M^0_{\rm env}}\right)\left(\frac{R_{\rm star}}{R_0}\right)^{-1},\\ C &= \left(\frac{M_2}{M_\odot}\right)\left(\frac{M_{\rm env}}{M^0_{\rm env}}\right)\left(\frac{R_{\rm star}}{R_0}\right)^{-1},\\ D &= A + \frac{M_2}{M_\odot}\,,\\ \tau_0^{-1} &= \frac{8\pi G^{1/2}\rho_0 r_0^{3/2}}{M_\odot^{1/2}}\,. \end{aligned} \tag{15}$$

For the above assumed parameters and a core mass of $M_c \sim$ 0.6 $M_\odot$, equation (13) gives a spiraling-in time scale (from $r \simeq 100\,R_\odot$ inward) of less than 20 yr. While this should defi-

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1994ApJ...421..219S

nitely be considered as a lower limit, since our formalism did not take into account neither the process of spin-up of the envelope, which slows down orbital decay (e.g., Livio & Soker 1988), nor the reduction in the density of the CE due to mass motions, this calculation does demonstrate that the evolution in the final stages of CE evolution is very rapid. Consequently, the entropy profile in the secondary star remains frozen during these stages (the mass-loss phase is essentially adiabatic), this was confirmed by the numerical calculations of Hjellming & Taam.

The direct outcome of the above sequence of events is that as the CE phase ends (with the ejection of the envelope), the secondary will expand adiabatically (responding to the new boundary conditions). Since at this stage the binary can be quite close, *mass transfer can ensue and an accretion disk will form.*

One may ask what is the difference between this mass transfer episode and mass transfer in say, cataclysmic variables (CVs), since in both cases mass is being transferred from a low-mass main-sequence star to a compact component. This is an important question, because in spite of the fact that disks are certainly present in CVs, no jets are directly observed in these systems (although wind outflows are certainly observed, e.g., Drew 1991). The question is therefore: is jet formation more likely in the post–CE systems described here?

There are two main differences between the mass transfer episode which follows the CE phase and the regular mass transfer process in CVs.

First, the mass transfer rate in the post–CE binary can be significantly higher than typical transfer rates in CVs. The disturbed surface layers of the secondary star typically involved a mass of $\Delta M \sim 0.01$–$0.05\ M_{\odot}$ in the calculations of Hjellming & Taam. Since the characteristic thermal relaxation time for these layers was of order 10^3–10^4 yr, mass transfer rates of order 10^{-6}–$5 \times 10^{-5}\ M_{\odot}\ \mathrm{yr}^{-1}$ can be expected. This is 10^2–10^5 times higher than characteristic transfer rates in CVs.

Second, the pre-white dwarf accreting object that is the central star of the PN is considerably hotter than white dwarfs in CVs.

The first fact makes a strong outflow very plausible, since in CVs the outflow rate is typically found to be of order 1%–10% of the mass flow rate through the disk (Drew 1991) and winds are observed only in high-mass transfer systems (nova like variables and dwarf novae in outburst). The fact that the accreting object is hot also facilitates mass loss. In fact, the hot central stars of PNs are expected to emit a hot dilute wind even by themselves (e.g., Kwok 1982; Kahn 1982). These two effects may explain why jets can be obtained more easily in post–CE central stars than in CVs.

Although we are unable to calculate the evolution from the CE phase up to the formation of jets, we can expect the jets (if they indeed form) to have the following characteristics. The mass-loss rate into the jets will be of the order of 10% of the mass-loss rate from the secondary, i.e., $\sim 10^{-6}\ M_{\odot}\ \mathrm{s}^{-1}$. This is a higher mass-loss rate than that in the cases discussed in the previous section and in Soker (1992a), but lower than in the case of a secondary outside the envelope. At this stage the primary radius is $\sim 1\ R_{\odot}$, and the jets velocities will be $\sim 500\ \mathrm{km\ s^{-1}}$. Even though most of the material is concentrated in the equatorial plane (because of the influence of the secondary during the CE phase) some matter will be present in the polar direction. The jets will propagate through this medium and interact with it. These jets might be similar in properties to HH objects, with the additional ionization source from the central star, and possibly, with a massive torus in the equatorial plane.

5. SUMMARY

We have studied three possible scenarios for the formation of disks, which may form jets, in PNs. In the first (§ 2), which is due to Morris (1987), a compact companion accretes mass from the progenitor AGB star, and a disk is formed around this companion. Bipolar nebulae which seem to fit this model have been actually observed (e.g., He2-104, BI Cru; Corradi & Schwarz 1993, and references therein). We quantified this model somewhat and found that these systems are likely to be concentrated toward the galactic plane. This process is qualitatively different from the other scenarios, in that here it is the companion star which is collimating the outflows.

In the second scenario (§ 3), the basic process is the accretion of backflowing nebular material by the CS. If there is a relative angular momentum between the backflowing gas and the CS (due, say, to the presence of a companion), then a disk may form. This process (if it occurs) can take place only in the very late stages of the PN phase, or even after the PN has faded. This scenario cannot explain the majority of jets, which are observed much before the PN started fading. However, two points should be noted here. First, the backflowing, hydrogen-rich material, can enrich the cooling WD with hydrogen. And second, a late helium flash may reveal the existence of these late jets. This scenario needs further study in relation to the last two points.

The third (and probably most likely) scenario (§ 4) deals with the exit from a CE phase, when the secondary is a main-sequence star. We noted that since the secondary is out of thermal equilibrium when it emerges from the CE, and it loses most of the matter it had accreted during the CE phase (Hjellming & Taam 1991), a disk around the primary is likely to form.

In addition to the configurations studied in the present work, we should mention the suggestion by Soker (1992a) that jets form by collimated outflows from the distorted envelope of the star as it leaves the AGB. The distortion itself was assumed to result from a low-mass secondary orbiting inside the CE.

In the present work, we emphasized the role of binaries in the generation of axially symmetric outflows. It is presently unclear whether or not single stars are also capable of producing such flows (see, e.g., Livio 1992 for a discussion). In particular, it is not known yet if the same mechanism which may produce equatorial density enhancements in Be stars (Bjorkman & Cassinelli 1993) can operate also in AGB stars. This process is currently under study.

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1994ApJ...421..219S

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