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APPLICATION OF MAGNETOHYDRODYNAMIC DISK WIND SOLUTIONS TO PLANETARY AND PROTOPLANETARY NEBULAE

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ABSTRACT

Winds from accretion disks have been proposed as the driving source for precessing jets and extreme bipolar morphologies in planetary nebulae (PNs) and proto-PNs (PPNs). Here we apply MHD disk wind models to PNs and PPNs by estimating separately the asymptotic MHD wind velocities and mass-loss rates. We discuss conditions that may occur in PN and PPN accretion disks that form via binary interactions. We show that the resulting winds can recover the observed momentum and energy input rates for PNs and PPNs. High accretion rates ($\dot{M}_a \approx 10^{-4}\ M_\odot\ \mathrm{yr}^{-1}$) may be required in the latter case. We find that the observed total energy and momentum in PPNs can be recovered with disk wind models using existing disk formation scenarios. When combined with existing scenarios for accretion disk formation from disrupted stellar companions, our models may provide an explanation for the existence of high-speed polar knots (FLIERS) observed in some PNs.

Subject headings: magnetic fields — MHD — planetary nebulae: general — shock waves

1. INTRODUCTION

Planetary nebulae (PNs) and proto–planetary nebulae (PPNs) are believed to be the penultimate evolutionary stages of low- and intermediate-mass stars ($M_* \leq 8\ M_\odot$). PNs and PPNs appear on the sky as expanding plasma clouds surrounding a luminous central star. As ground-based telescopes increased their resolution, elliptical and bipolar PNs were revealed (Balick 1987). The bipolar nebulae may be further subdivided into "butterfly" shapes in which the "waist" is pinched at the central star and "bilobed" PNs in which a pair of larger outer lobes connects to a central and generally smaller round or elliptical nebula (for a review see Balick & Frank 2002 and references therein). More recently, deeper and higher resolution studies with the *Hubble Space Telescope* have shown evidence of narrow collimated features that appear to be better described as "jets" than "bipolar lobes." The nature of these jets and other highly collimated bipolar outflows in both PNs and PPNs remains a subject of considerable debate (Sahai & Trauger 1998).

Considerable progress has been made in understanding the hydrodynamic shaping of elliptical and bilobed PNs (for reviews, see Frank 1999; Balick & Frank 2002); however, the origin of extreme butterfly nebulae as well as jets in PNs continues to pose a number of problems for theorists. The outstanding issues include (1) the lack of a large-scale torus to focus the outflows, (2) the very large velocities in some jets, and (3) the point symmetry of outflows.

Beyond issues of morphology, there exists a formidable problem for PPNs concerning the total momentum and energy in the outflows. A number of observational studies have shown that radiatively accelerated winds in PPNs cannot account for the high momentum and energy implied by CO profiles. This problem was identified first by Knapp (1986) and was most recently and comprehensively investigated by Bujarrabal, Alcolea, and collaborators (Bujarrabal et al. 2001). The latter find that 21 of 23 CO-emitting PPN objects show outflows whose scalar momenta ($\Pi = MV$) are more than 10^3 times larger than those in the stellar radiation. Thus, $\Pi \gg (L_*/c)\Delta t$, where L_* is the stellar luminosity emitted during the PPN outflow expansion lifetime Δt. In light of these results, both the launching and the collimation of winds in PPNs becomes problematic.

The dominant hydrodynamic theory for shaping PNs had been the generalized interacting stellar winds (GISW) model (Kwok, Purton, & Fitzgerald 1978; Balick 1987; Icke 1988) in which a star and its wind evolve from the asymptotic giant branch (AGB) to a white dwarf. A slow, dense wind expelled during the AGB phase is followed by a fast, tenuous wind driven off the contracting proto–white dwarf during the PN phase. Numerical models have shown that this paradigm can embrace a wide variety of nebular morphologies including highly collimated jets (Icke et al. 1992; Mellema & Frank 1997; Borkowski, Blondin, & Harrington 1997) when the slow wind takes on aspherical density distributions. While the GISW model *can* produce narrow jets, it usually requires a large-scale "fat" torus. It is difficult to imagine that a large-scale outflowing gaseous torus can provide a stiff precessing nozzle for production of point-symmetric flows. (Recent results of Icke 2003, however, indicate that pure hydrodynamics may lead to point-symmetric shapes in some cases.)

In addition, it is now recognized that fast ($\geq$100 km s^{-1}) bipolar outflows can occur in the PPN or even in the post-AGB stage. Objects like CRL 2688 and OH 231.8+4.2 raise the question of how high-velocity collimated flows occur when the star is still in a cool giant or even supergiant stage (CRL 2688 has an F supergiant spectral type). Finally, the GISW model assumes a radiation-driven wind. As discussed above, the results of Bujarrabal et al. (2001) make radiation driving untenable as a source for many PPNs flows.

Models invoking a toroidal magnetic field embedded in a normal radiation-driven stellar wind have shown considerable promise (Luo & Chevalier 1994; Różyczka & Franco 1996; García-Segura 1997; García-Segura et al. 1999). Recent results, however, also imply that jets may form at smaller distances and that such models may not begin with appropriate initial

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conditions (Gardiner & Frank 2001). In addition, by their very nature, magnetized wind bubble models cannot account for the momentum excesses in PPNs since they also require radiation-driven winds. The fields are simply too weak to power the observed outflows.

Thus there remains considerable uncertainty about the processes that produce collimated jets in PPNs and PNs. Other systems that produce jets, such as young stellar objects (YSOs), active galactic nuclei (AGNs), and microquasars, have been modeled via a combination of magnetic and centrifugal forces from accretion disks (Blandford & Payne 1982; Shu et al. 1994; Königl & Pudritz 2000, hereafter KP00).

The success of these ''magnetocentrifugal launching'' (MCL) models is such that it is worthwhile considering whether such a scenario can be applied to PNs and PPNs. Indeed, Morris (1987) and Soker & Livio (1994) mapped out scenarios in which accretion disks form around binary PN progenitors. Each study equated the existence of disks with the existence of jets. The details of the jet launching and collimation mechanism where not, however, specified. Recent work by Soker & Rapport (2001), Soker (2001), and Livio & Soker (2001) have relied heavily on collimated winds from disks, but these works also do not specify how such winds are launched or collimated. Thus, application of MCL disk wind models to PNs and PPNs would close an important gap in building a new paradigm for these systems. In particular, MCL models may offer a means of resolving both issues associated with PN jet precession (the underlying disk precesses via instabilities; Livio & Pringle 1996; Quillen 2001) *and* those associated with PPN momentum excesses (no need for radiation pressure launching winds).

Recent studies by Blackman et al. (2001a) and Blackman, Frank, & Welch (2001b) have explored MCL paradigms for both the star and disk in PPN and PN systems. They computed an upper limit on the magnetic luminosity available to power an outflow, assuming that a dynamo is the source of the large-scale magnetic field. In this paper we provide further calculations along these lines, deriving scaling relations from the equations for MCL, and separately estimate the mass outflow rate and the asymptotic outflow velocity. We then compare these results with PNs and recent observations of PNs. When the dynamo is invoked to produce the field, the mechanical wind luminosity and thus outflow rate are naturally linked to the accretion rate.

In § 2 we derive the outflow speed and the mass-loss rate by combining results from magnetically driven wind theory and dynamos. In § 3 we discuss models of PPN and PN disks. In § 4 we apply the results of §§ 2 and 3 to PNs and PPNs and show that it is easy to solve the aforementioned PPN momentum and energy excess problems. In § 5 we conclude and discuss open questions.

2. MASS OUTFLOW RATE AND ASYMPTOTIC WIND SPEED FROM MCL THEORY AND DYNAMO THEORY

2.1. *Magnetic Luminosity*

The basic physics of magnetocentrifugal launching of winds and jets is well studied when a magnetic field distribution is imposed on the disk (Blandford & Payne 1982; Sakurai 1985; Pelletier & Pudritz 1992; Shu et al. 1994; Ostriker 1997; KP00). These models presume an initial field of a given strength and geometry, but the tendency for large-scale fields to diffuse (Lubow, Papaloizou, & Pringle 1994; Blackman 2004) suggests that the field must be generated in situ by a dynamo (e.g., Blackman & Field 1999). In this section we briefly summarize how to combine the basics of Poynting flux–driven outflows with asymptotic wind solutions and mean field dynamo theory to estimate the asymptotic wind speed and the outflow accretion rate.

Magnetocentrifugal launching is a means of converting gravitational binding energy in an accreting source into kinetic energy of an outflowing wind. The magnetic fields act as a drive belt to extract angular momentum from the anchoring rotator and launch the wind. The magnetic luminosity, or equivalently, the maximum magnetic power available for a wind, can be obtained from the integrated Poynting flux (Blackman et al. 2001b). The field lines rotate nearly rigidly at angular speeds associated with the anchoring footpoint $\Omega_0(r)$ up to the Alfvén radius $r_{\rm A}(r)$. After this point the angular speed falls off with $1/r^2$ (conserving specific angular momentum) and the field falls off as $1/r$. The field is primarily poloidal out to the Alfvén radius associated with each field line, where the poloidal and toroidal components are comparable. If the poloidal field falls off as $1/r^2$ out to the Alfvén radius and is not too far from the disk surface, the Poynting flux can be approximated by the contribution from the Alfvén radius associated with field lines anchored at the inner radius of the disk. That is, we have

$$L_w = \frac{1}{2}\dot{M}_w u_\infty^2 \sim L_{\rm mag} \equiv \int (\boldsymbol{E}\times\boldsymbol{B})\cdot d\boldsymbol{S}_{\rm A} \sim \int_{r_i}^{r_{\rm A}(r_i)} \Omega(r) r B_p B_\phi r\, dr \sim B_{\rm A}^2 \Omega_0 r_{\rm A}^3, \quad (1)$$

where r_0 is the disk inner radius and where $B_{\rm A} = B_\phi \sim B_p$ at the Alfvén surface (i.e., the toroidal and poloidal field components are nearly equal). Since the dominant contribution to the magnetic luminosity comes from the field lines anchored at the innermost radius r_i, in what follows, all quantities labeled with subscript 0 refer to those values evaluated at $r = r_i$, the innermost disk radius. This also implies that $r_0 = r_i$.

2.2. *The Bernouilli Constant*

To estimate the mass outflow rate and outflow speed separately, a bit more work is required. The MCL problem requires the construction of solutions for a steady, ideal, isothermal magnetohydrodynamic flow. The isothermal assumption eliminates the need for solving the energy equation, but more complex assumptions can be used, e.g., a polytropic law. The system of equations to solve become mass conservation, momentum conservation, and the steady state induction equation (e.g., Pelletier & Pudritz 1992). Respectively, these are

$$\nabla\cdot(\rho\boldsymbol{u}) = 0, \quad (2)$$

$$\rho\boldsymbol{u}\cdot\nabla\boldsymbol{u} = \frac{1}{4\pi}(\nabla\times\boldsymbol{B})\times\boldsymbol{B} - \nabla P - \rho\nabla\Phi, \quad (3)$$

$$\nabla\times(\boldsymbol{u}\times\boldsymbol{B}) = 0, \quad (4)$$

where Φ is the gravitational potential due to the central source. In cylindrically symmetric coordinates, the physics of

MCL disk winds can be cast in axisymmetric form where the velocity and magnetic fields are decomposed into toroidal (ϕ) and poloidal (r, z) components,

$$\boldsymbol{B} = \boldsymbol{B}_p + B_\phi \hat{\boldsymbol{\phi}}, \tag{5}$$

$$\boldsymbol{u} = \boldsymbol{u}_p + \boldsymbol{u}_\phi = \boldsymbol{u}_p + \Omega r \hat{\boldsymbol{\phi}}, \tag{6}$$

where Ω is the rotational frequency of the plasma at a point (r, z). Axisymmetry and $\nabla \cdot \boldsymbol{B} = 0$ allow $\boldsymbol{B}_p$ to be expressed in terms of a magnetic flux function $a(r, z)$ such that

$$\boldsymbol{B}_p = \frac{1}{r}(\nabla a \times \hat{\boldsymbol{\phi}}). \tag{7}$$

The magnetic surfaces, generated by rotation of a poloidal field line about the axis, are surfaces of constant $a(r, z)$.

For axially symmetric configurations, equations (2) and (4) imply that the poloidal velocity is always parallel to magnetic surfaces, that is,

$$\rho \boldsymbol{u}_p = k(a) \boldsymbol{B}_p, \tag{8}$$

where the function k is constant on a flux surface (KP00, eq. [7]). The induction equation and $\nabla \cdot \boldsymbol{B} = 0$ also imply

$$\rho u_\phi = k(a) B_\phi + \rho r \Omega_0(a), \tag{9}$$

where $\Omega_0(a)$ is a constant that is approximately equal to the angular velocity in the disk where the magnetic surface is tied (KP00, eq. [9]).

Using $\nabla \cdot \boldsymbol{B} = 0$, equation (2), and the azimuthal component of equation (3), the conserved angular momentum per unit mass $l(a)$ becomes

$$l(a) = \Omega r^2 - \frac{r B_\phi}{4\pi k} = \text{const}\,(a) = \Omega_0 r_\text{A}^2, \tag{10}$$

where the last equality follows from using equation (9) and finding the value at the Alfvén radius $r_\text{A}(a)$ for a magnetic flux surface anchored to the disk at radius $r_0(a)$ (e.g., KP00, eq. [11]). The Alfvén radius $r_\text{A}(a)$ defines the radial coordinate of the point along a poloidal magnetic surface when the outflow speed on that surface equals the local poloidal Alfvén velocity $u_\text{A} = B_p/(4\pi\rho)^{1/2}$. The expression for conservation of angular momentum (eq. [10]) thus relates contributions from the angular momentum of matter and magnetic torques to a fiducial value associated with the Alfvén point.

The poloidal component of equation (3) can be integrated using equations (9) and (10) to give the total conserved specific energy $U(a)$ carried by the wind on magnetic surfaces in both kinetic energy and Poynting flux. By defining w as the enthalpy per unit mass, the generalized Bernoulli integral then emerges (KP00, eq. [14]) as

$$\frac{1}{2}\left(u_p^2 + \Omega^2 r^2\right) + \Phi + w + \Omega_0\left(\Omega_0 r_\text{A}^2 - \Omega r^2\right) = U(a) = \text{const}\,(a). \tag{11}$$

2.3. *The Outflow Speed and Mass Outflow Rate*

We now assume a cold wind (such that w can be ignored). Because little acceleration occurs outside the Alfvén surface, we assume that

$$u_\infty = f u_\text{A}, \tag{12}$$

where $f \gtrsim 1$. We solve for u_A and constrain f below.

To calculate u_A, we now solve the momentum equation for $r \leq r_\text{A}$ in the rotating frame. For a cold wind, the dominant force components are the centrifugal force and the gravitational force. For $r < r_\text{A}$, the magnetic field lines can be assumed to rotate rigidly with the angular speed of their footpoints at $r = r_0$. It is straightforward to see that the steady state radial momentum equation can then be written (Blandford & Payne 1982)

$$\boldsymbol{u} \cdot \nabla u_r = -\partial_r \Phi_\text{eff} = \left(\frac{GM_*}{r_0}\right)\left[\frac{r}{r_0^2} - \frac{r r_0}{(z^2 + r^2)^{3/2}}\right]. \tag{13}$$

To keep the analysis as simple as possible, consider the initial launch to be highly inclined, such that $z \ll r$. Such an approximation is reasonable since equation (13) implies that the field inclination from the disk must make an angle of less than $\pi/4$ to the disk plane for launch. This result for the field inclination follows from equation (13) by considering $z \ll r$ and $(r - r_0) \ll r_0$ and carrying out an expansion of the resulting force along the field line to second order in z^2/r_0^2 and $(r - r_0)^2/r_0^2$. The reason for an outward force is the gain in centrifugal force at $r > r_0$ from the fact that "rigid" field lines enforce corotation. (Note that in the following sections we use the fact that the dominant contribution to the magnetic luminosity comes from the field lines anchored at the innermost radius r_i, implying that $r_0 = r_i$.)

Integrating equation (13) along the field line for the case $z \ll r$ and taking the result at $r = r_\text{A}$ gives

$$u_\text{A}^2 \simeq \Omega_0^2 r_\text{A}^2 \left(1 + 2\frac{r_0^3}{r_\text{A}^3}\right) = q^2 \Omega_0^2 r_\text{A}^2, \tag{14}$$

where we see that

$$1 < q^2 = \left(1 + 2\frac{r_0^3}{r_\text{A}^3}\right) < 3. \tag{15}$$

Now

$$u_\text{A}^2 = \frac{B_\text{A}^2}{4\pi\rho_\text{A}}, \tag{16}$$

where ρ_A is the mass density at r_A and

$$\rho_\text{A} = \frac{\dot{M}_{w,l}/\xi}{r_\text{A}^2 u_\text{A}} = \frac{\dot{M}_w}{4\pi r_\text{A}^3 \Omega_0 q} \tag{17}$$

(from the mass continuity equation), where ξ is the solid angle $\leq 4\pi$ corresponding to the launched outflow mass-loss rate $\dot{M}_{w,l}$, whereas $\dot{M}_w = 4\pi \dot{M}_{w,l}/\xi$ is the effective mass-loss rate were it quasi-spherical.

Using equations (14), (16), and (17), we then obtain

$$\dot{M}_w = \frac{B_A^2 r_A}{q\Omega_0}. \quad (18)$$

But using equations (12) and (1), we get a separate equation for $\dot{M}_w$, namely,

$$\dot{M}_w = \frac{2}{f^2 q^2} \frac{B_A^2 r_A}{\Omega_0}. \quad (19)$$

By setting equation (18) equal to equation (19), we obtain $f^2 q = 2$. Thus from equation (15) we must have

$$\frac{2}{3^{1/2}} < f^2 < 2. \quad (20)$$

This is a narrow range of f, and for our crude order-of-magnitude estimates we will take $f = 1.2$, a value right in the middle of the allowed range. For this choice of f, we obtain $q = 1.4$. For simplicity, we use these values in what follows. From equation (15) these values imply

$$r_A = 1.27 r_0. \quad (21)$$

Using these in equations (12) and (18) then gives

$$u_\infty \simeq 2.1 \Omega_0 r_0, \quad (22)$$

$$\dot{M}_w \simeq 0.9 \frac{B_A^2 r_0}{\Omega}. \quad (23)$$

In the next subsection we obtain an expression for B_A.

2.4. *Magnetic Field Strength*

Magnetic fields may form in these disks via dynamo processes (Reyes-Ruiz & Stepinksi 1995; Blackman et al. 2001b). The topology of such a field (the ratio of poloidal B_p and toroidal B_ϕ in the disk) and its subsequent value in the coronae remain a subject of considerable discourse (e.g., Blackman 2004). Here we assume that a field produced by dynamos can drive a disk wind in the manner described in the last section. This means that whatever combination of toroidal and poloidal field is produced in the disk, we assume a primarily poloidal field in the corona where the wind launches.

Equation (4) and $\nabla \cdot \boldsymbol{B} = 0$ imply that the radial magnetic field falls off with r^2 along the field line to r_A. Thus,

$$B_A^2 = B_0^2 \left(\frac{r_0}{r_A}\right)^4 = 0.4 B_0^2, \quad (24)$$

where B_0 is the poloidal field at the disk surface and we have used equation (21). The square of the surface poloidal field B_0^2 can be estimated to be lower than the midplane poloidal field squared, B_d^2, by the density ratio to the 4/3 power (flux freezing). The density falls by a factor of e^{-n}, where n is the number of scale heights above the disk from where the wind is launched. We then have

$$B_0^2 = e^{-4n/3} B_d^2 \sim 4\pi e^{-4n/3} \rho_d \alpha_{SS}^2 c_s^2, \quad (25)$$

where α_{SS} is the disk viscosity parameter (Shakura & Sunyaev 1973), c_s is the sound speed, and ρ_d is the disk density. The latter similarity in equation (25) follows from estimating the disk mean poloidal field from a helical mean field dynamo (Blackman et al. 2001b; Blackman 2004). From mass conservation in the disk, the disk density satisfies

$$\rho_d = \frac{\dot{M}_d}{2\pi r_0 h_0 v_r}, \quad (26)$$

where h_0 is the disk scale height at $r = r_0$, $v_r = \alpha_{SS} c_s h_0 / r_0$ is the disk radial accretion velocity, and $\dot{M}_d$ is the mass accretion rate. Combining this with equations (24), (25), and (26) then implies

$$B_0^2 = \frac{2e^{-4n/3} \alpha_{SS} c_s \dot{M}_d}{h_0^2} = 2.5 B_A^2. \quad (27)$$

2.5. *Compiling the Formulae for Application to Observations*

Combining equation (27) with equation (23) and using $c_s = \Omega h_0$ for an accretion disk, we obtain

$$\dot{M}_w = 0.72 e^{-4n/3} \alpha_{SS} \frac{r_0}{h_0} \dot{M}_d \quad (28)$$

for the mass-loss rate. Combining this with equations (22) and (1) gives

$$L_{mag} \simeq 1.6 e^{-4n/3} \alpha_{SS} \frac{r_0}{h_0} \dot{M}_d \Omega_0^2 r_0^2. \quad (29)$$

The scaling for the momentum input $\dot{\Pi} = \dot{M}_w u_\infty$ of the wind can then also be estimated via

$$\dot{\Pi} \sim \frac{L_{mag}}{u_\infty} \simeq 0.76 e^{-4n/3} \alpha_{SS} \frac{r_0}{h_0} \dot{M}_d \Omega_0 r_0. \quad (30)$$

In the above three relations, there is some freedom in choosing n, the number of scale heights above the midplane from which the wind launches (i.e., the location at which the plasma becomes magnetically dominated). One expects the corona to become magnetically dominated after 1 or 2 scale heights; thus $n = 1.5$ is reasonable choice for the point at which we expect the ratio of thermal to magnetic pressure (β) to make the transition $\beta > 1$ to $\beta < 1$. Using this we obtain

$$\dot{M}_w \sim 0.1 \alpha_{SS} \frac{r_0}{h_0} \dot{M}_d, \quad (31)$$

$$L_{mag} \simeq 0.22 \alpha_{SS} \frac{r_0}{h_0} \dot{M}_d \Omega_0^2 r_0^2, \quad (32)$$

$$\dot{\Pi} \sim \frac{L_{mag}}{u_\infty} \simeq 0.14 \alpha_{SS} \frac{r_0}{h_0} \dot{M}_d \Omega_0 r_0. \quad (33)$$

The last two expressions give estimates of the rate that energy and momentum are input by MCL disk winds to the ambient medium, and equation (22) gives the asymptotic wind velocity

$$u_\infty \sim 2.1 \Omega_0 r_0. \quad (34)$$

These can be used for comparison with observations.

3. DISK ACCRETION RATE IN PNs

In order to produce a more detailed comparison of MHD disk winds with PNs, it is necessary to have a model for PN accretion disks. In particular, it is necessary to know the accretion rate $\dot{M}_d$ and the inner disk radius r_i as a function of time.

It is unlikely that an accretion disk could survive the long main-sequence lifetime of a PN central star. Thus, unlike YSOs and AGNs, accretion disks in PN systems must form via binary interactions. Disks may form around secondaries via Roche lobe overflow or accretion of the dense AGB wind (Morris 1987; Mastrodemos & Morris 1998). Such systems would be similar to symbiotic stars (Corradi et al. 2000). Accretion disks could also form around the primary after common-envelope (CE) evolution and disruption of the secondary star (Soker & Livio 1994; Soker 1998; Reyes-Ruiz & López 1999).

Mastrodemos & Morris (1998) carried out detailed simulations of the first scenario. Using a smoothed particle hydrodynamics (SPH) method, they modeled the gravitational interaction of a dense AGB wind with a lower mass companion. They found steady accretion disks around a white dwarf companion orbiting an AGB star with $\dot{M}_{\rm AGB} \approx 10^{-5}\ M_\odot\ {\rm yr}^{-1}$. The ratio of $\dot{M}_d/\dot{M}_{\rm AGB} \approx 0.05-0.005$ that they found in their models is consistent with expectations from basic theory (Frank, King, & Raine 2002).

Accretion disks may also form via disruption of the secondary after CE evolution (Soker & Livio 1994). This model implies a finite lifetime for the disk as the mass reservoir of the disrupted companion is slowly drained onto the primary. A description of disk formation in PNs has been given in Reyes-Ruiz & López (1999). Envelope ejection occurs via transfer of angular momentum during which the secondary falls to a separation such that it catastrophically overflows its Roche lobe and forms a disk around the primary. There are a number of important constraints on the properties of binaries that would lead to disk formation in this way. Reyes-Ruiz & López (1999) found that systems with a primary consisting of an evolved AGB star with mass $M_* \approx 2.6-3.6\ M_\odot$, a low-mass secondary ($\lesssim 0.08\ M_\odot$), and an initial binary separation of less than 200 $R_\odot$ may produce disks. The AGB star will shed most of its mass during the CE ejection, leaving a post-AGB stellar core surrounded by a thin shell.

Reyes-Ruiz & López (1999) find the disk accretion rate to evolve in time in a power-law manner:

$$\dot{M}_d = \dot{M}_{d0}\left(\frac{t}{1\ {\rm yr}}\right)^{-5/4} M_\odot\ {\rm yr}^{-1}. \quad (35)$$

A typical value of the scale is $\dot{M}_{d0} = 10^{-3}\ M_\odot\ {\rm yr}^{-1}$.

4. DISK WINDS MODELS FOR PNs AND PPNs

We wish to understand whether disk wind models can account for outflows in PNs and PPNs and to characterize the parameters for the winds. The momentum $\dot{\Pi} = \dot{M}_w u_w$ and energy $\dot{E} = \frac{1}{2}\dot{M}_w u_w^2$ injection rates for PNs are easily approximated. For "classic" PNs, a total mass of $M_{\rm PN} \approx 0.1\ M_\odot$ must be accelerated to a velocity of $u_{\rm PN} \approx 40$ km s^{-1} on a timescale of order $\Delta t_{\rm PN} \approx 10{,}000$ yr. (Note that the subscript PN refers to observed properties of the total nebula. This reflects material swept up by the wind modeled in this paper.) This gives $\dot{\Pi} = M_{\rm PN}u_{\rm PN}/\Delta t_{\rm PN} \approx 10^{27}$ g cm s^{-2} and $\dot{E} = M_{\rm PN}u_{\rm PN}^2/\Delta t_{\rm PN} \approx 10^{34}$ ergs s^{-1}.

Recall that Bujarrabal et al. (2001) found high total outflow momentum $10^{36} < \Pi/({\rm g\ cm\ s^{-1}}) < 10^{40}$ and total outflow energy $10^{41} < E/({\rm ergs\ s^{-1}}) < 10^{47}$ in an extensive sample of PPNs. These values can be converted into momentum and energy injection rates using an assumed injection or "acceleration" timescale Δt: $\dot{\Pi} = \Pi/\Delta t$ and $L = E/\Delta t$. Recall that the values for Π and E quoted above cannot be explained via radiation wind driving. Note also that there is uncertainty about Δt, the injection timescale, but most likely $\Delta t < 10^3$ yr. The question that then arises is, can energy and momentum budgets be met with disk wind models?

The results of § 2 demonstrate that disk wind solutions are sensitive to the location of footpoints for the flow r_0 and the accretion rate $\dot{M}_d$. In what follows we assume that the inner edge of the disk extends to the stellar surface and use $r_0 = r_i = r_*$. We shall see that the accretion rate is the parameter that becomes most important for obtaining solutions for PPNs since the requisite high-outflow momenta will require high values of $\dot{M}_d$. One means of achieving high accretion rates will be to use the model described by Reyes-Ruiz & López (1999), where accretion rates as high as a few times $10^{-4}\ M_\odot\ {\rm yr}^{-1}$ are possible for short periods as the disrupted companion's mass is fed onto the surface of the primary.

PN solutions.—We first consider the case of a classic PN. In this case we would consider that the star that produces the jet is a proto–white dwarf with an AGB companion (Soker & Rappaport 2000). Thus accretion rates of $\dot{M}_d \approx 10^{-6}$ are reasonable. To evaluate the expressions above we choose parameters for a canonical PN with mass $M_s = 0.6\ M_\odot$ and a disk with $\alpha = 0.1$ and $r_0/h_0 = 10$.

If we assume typical PN central star parameters ($T_* = 10^5$ K and $L_* = 5000\ L_\odot$ such that $r_i = 1.64 \times 10^{10}$ cm), we find the following conditions for the wind from equations (31)–(34) (note that we use $u_w = u_\infty$):

$$\dot{M}_w = 1 \times 10^{-7}\ M_\odot\ {\rm yr}^{-1}\left(\frac{\dot{M}_a}{10^{-6}\ M_\odot\ {\rm yr}^{-1}}\right), \quad (36)$$

$$u_w = 1.25 \times 10^3\ {\rm km\ s}^{-1}\left(\frac{M_*}{0.6\ M_\odot}\right)^{1/2}\left(\frac{r_0}{0.23\ R_\odot}\right)^{-1/2}. \quad (37)$$

Thus using typical conditions for PN central stars, the scaling relations derived from the MHD equations yield disk wind parameters well matched with observations.

PPN solutions.—While the mass-loss rates and velocities are known for PN winds, the situation for PPNs is not as clear. In general, what is observed in PPNs is the total mass in the outflows. Mass-loss rates must be inferred from the estimates of PPN acceleration timescales. Velocities are also uncertain in the sense that the winds themselves cannot be observed directly but only properties of swept-up material can be determined.

We assume a model post-AGB star with a mass of $M_s = 0.6\ M_\odot$, a temperature of $T_* = 10{,}000$ K, and a luminosity of $L_* = 5 \times 10^3\ L_\odot$, which, assuming a blackbody, yields a radius of $r_* = 1.6 \times 10^{12}$ cm $= 23\ R_\odot$. Note that such a star has an escape velocity of $u_{\rm esc} = 98$ km s^{-1}.

From our previous discussion it is clear that achieving the high momentum input rates observed in PPNs via MCL disk wind models will necessitate high accretion rates. Thus we adopt an acceleration timescale of $\Delta t = 200$ yr and an accretion rate of $\dot{M}_d = 1 \times 10^{-4}\ M_\odot\ {\rm yr}^{-1}$. This value of $\dot{M}_d$ is the 200 yr average of that found by Reyes-Ruiz & López (1999) for their case A ($\dot{M}_{d0} = 1.6 \times 10^{-3}\ M_\odot\ {\rm yr}^{-1}$; eq. [35]). Once again we choose a disk with $\alpha = 0.1$ and $r_0/h_0 = 10$.

Assuming $r_0 = r_*$ along with the other parameter values given above, the key wind quantities $\dot{M}_w$, $u_w = u_\infty$, $L_m = L_w$, and $\dot{E}$ can again be determined from equations (31), (34), and (32):

$$\dot{M}_w = 1. \times 10^{-5}\ M_\odot\ \mathrm{yr}^{-1}\left(\frac{\dot{M}_d}{10^{-4}\ M_\odot\ \mathrm{yr}^{-1}}\right), \tag{38}$$

$$u_w \simeq 146\ \mathrm{km\ s}^{-1}\left(\frac{M_*}{0.6\ M_\odot}\right)^{1/2}\left(\frac{R_i}{23\ R_\odot}\right)^{-1/2}, \tag{39}$$

$$L_m \simeq 6.7 \times 10^{34}\ \mathrm{ergs\ s}^{-1}\left(\frac{\dot{M}_d}{10^{-4}\ M_\odot\ \mathrm{yr}^{-1}}\right) \times \left(\frac{M_*}{0.6\ M_\odot}\right)\left(\frac{R_i}{23\ R_\odot}\right)^{-1}. \tag{40}$$

Note that $L_w \Delta t \approx 10^{44}$, a value in the middle of the range found by Burharrabal et al. (2001). Note also that the solution above has $u_\infty \approx 1.5 u_{\rm esc}$. Since $u_\infty \approx u_{\rm esc}$, the higher velocity outflows seen in some PPNs would require more disks around more compact central sources. We note here that the observed momenta and energy in the outflows come primarily from swept-up circumstellar material. Thus, as is the case in YSO molecular outflows (Bachiller 1996), the energy and momentum budget of the disk-driven outflow must be sufficient to power the observed outflows via so-called prompt or shock-driven entrainment.

Given a model of the temporal history of the disk accretion, the total energy and momentum for the outflows can be found. Replacing $\dot{M}_d$ in equation (35) with $\dot{M}_d(t)$ from the relations derived by Reyes-Ruiz & López (1999) and integrating with respect to t gives

$$E = \int \frac{1}{2}\dot{M}_w u_\infty^2 dt \approx 1.3 \times 10^{44}\ \mathrm{ergs}\left[1 - \left(\frac{1\ \mathrm{yr}}{t}\right)^{1/4}\right], \tag{41}$$

$$\Pi = \int \dot{M}_w u_\infty dt \approx 1.8 \times 10^{37}\ \mathrm{g\ cm\ s}^{-1}\left[1 - \left(\frac{1\ \mathrm{yr}}{t}\right)^{1/4}\right]. \tag{42}$$

These results show that the MCL disk wind models can achieve both energy and momentum injection rates *as well as* the total energy and momentum required to account for many PPNs described by Bujarrabal et al. (2001). The total energy and momenta budgets that we find from these solutions fall well within the range of PPN outflows with momentum excesses. Taken together with our previous calculations for classic PN winds, these results confirm the predictions of Blackman et al. (2001b) that magnetized disk winds can account for much of the outflow phenomena associated with collimated outflows in the late stage of stellar evolution.

The results above indicate that collimated flows that form from transient disks in the PPN stage will appear as dense knots in mature PN flows. This may also serve to explain the presence of so-called FLIERS (fast low-ionization emission regions; Balick et al. 1994) seen in some PNs. The mass-loss rates in the winds derived above rapidly decrease with time. Thus the bulk of the jet's mass will lie near its head. As the material in the disk is accreted onto the star, the jet will eventually shut off, leaving the dense knot to continue its propagation through the surrounding slow wind.

When the star makes its transition to a hot central star of a PN, its fast, tenuous spherical wind sweeps up a shell of the slow AGB wind material. The shell's expansion speed will typically be of order 40 km s^{-1}, and it will not catch up to the head of the jet. Thus during the PN phase the jet head will appear as a dense, fast-moving knot that should lie outside the PN's wind-blown bubble. We note that masses of FLIERs are estimated to be of order 10^{-4} to $10^{-5}\ M_\odot$, which is reasonable for the models presented above. FLIER velocities can be lower than the $\approx$100 km s^{-1} calculated above, but deceleration of the jet head will occur via interaction with the environment. We note also that hydrodynamic simulations of PN jets in which the jet ram pressure decreases in time (as would occur for our model) show characteristic patterns of backward-pointing bow shocks (apex pointing back toward the star; Steffan & López 1998). If such results are robust, the jets produced by disk winds in our scenario above may also yield similar morphologies.

5. DISCUSSION AND CONCLUSIONS

While purely hydrodynamic models for PN shaping provide an adequate description of many large-scale features seen in purely morphological studies, jets and hourglass-shaped bipolar nebulae appear to strain their explanatory power. In addition, the momentum and energy associated with many PPNs appear to be orders of magnitude larger than can be accounted for with radiation driving from the central star (even when multiple scattering is taken into account). In this paper we have derived and applied scaling relations derived from time-independent axisymmetric MHD equations to the winds and wind-driven outflows in PNs and PPNs.

For classic PNs (a central star with $T_* \gg 3 \times 10^4$ K), we find that magnetocentrifugal winds can account for typical observed wind properties. We find that only a 1% efficiency of accretion of an AGB companion wind is required to produce reasonable results.

Our results for PPNs show that momentum excesses need not occur for outflows driven by MCL winds. While this is encouraging in terms of finding a mechanism for driving PPN outflows, the solutions require fairly high accretion rates ($>10^{-5}\ M_\odot\ \mathrm{yr}^{-1}$). It is not clear if such conditions can be achieved with the frequency required by observations. While the solutions of Reyes-Ruiz & López (1999) yield accretion rates and time dependencies that lead to the correct outflow momenta and energetics, their models place fairly stringent limitations on the nature of the binaries that form disks from disrupted companions. If accretion onto undetected compact orbiting companions is invoked (Soker & Rappaport 2001), then higher values of $\dot{M}_d$ may not be required.

It is worth noting that these models imply an outflow from both the AGB wind and the jet. The AGB wind can, in principle, be sculpted by its own magnetic forces or can be swept at later times when a radiation-driven wind from the exposed core is initiated. Thus this class of models implies the possible misalignment between the jet and the main body of the nebula. In Blackman et al. (2001b) a model for such multipolar outflows was proposed in which magnetically launched outflows are driven from the both the disk and a nonaligned rapidly rotating AGB core. In this paper we did not consider the AGB wind to be magnetized and considered the field in the disk to arise via a turbulent dynamo there. If the AGB wind is magnetized, then shaping could occur with an alignment that is uncorrelated with that of the collimated disk wind. We note that current MHD models of AGB winds, however, tend to support the development of a torus (Matt et al. 2000) rather than a jet. Jet formation via initially weak fields in post-AGB winds have been explored (García-Segura et al. 1999); however, these models require higher winds speeds than occur on the AGB. Finally, it is worth noting that recent simulation results by Matt, Blackman, & Frank (2004) confirm that the exposed rapidly rotating magnetic core model can produce well-collimated outflows. Regardless of these caveats, the point remains that collimated disk winds and AGB wind systems can, in principle, explain multipolar outflows.

We note that the presence of molecules in fast-moving outflows driven by winds, as proposed in our studies, is not new. Studies of YSO molecular outflows have revealed "fast molecular" material moving at speeds of more than 25 km s^{-1}, which is the nominal dissociation speed for H_2 in J-shocks. A number of theoretical proposals to resolve the issue currently exist, including magnetic precursors in J-shocks (Harigan, Curiel, & Raymond 1989), high-velocity C-shocks (Smith & Brand 1990), accelerating shocks (Lim 2003), and the presence of clumped or inhomogeneous gas (Hartquist & Dyson 1987). Thus the behavior we see in PPNs also exists in YSOs, where compelling evidence for MHD-driven disk winds already exists. The question of how molecules can survive after being swept-up in high-velocity flows remains unanswered for PPNs and YSOs but is currently an active area of research with a number of competing models under investigation.

We note that a robust prediction of our models is the ratio of wind mass-loss rate to accretion rate, i.e., $\dot{M}_w/\dot{M}_a \approx 0.1$. This is true for most MCL disk wind models and can be seen as a target prediction that can be explored observationally.

Thus the question that must now be addressed is, can conditions for either high accretion rate disks in PPNs or symbiotic-type accreting companions be made to embrace enough systems to account for the statistics of Bujarrabal et al. (2001)? There are many uncertainties concerning the formation of disks via the disruption of the companion after a CE phase. The disruption of a secondary may not really lead to an "accretion disk" around the primary because of both tidal and hydrodynamic disruption. The requirement that material from the disrupted companion "push" through the envelope may not lead to a disk but rather to a dense and clumpy expanding torus. A detailed discussion of the process, however, remains outside the scope of our paper. Readers wishing to consider the viability of these models are encouraged to review Soker & Livio (1994) and Ruiz-Reyes & López (1999).

There is an important difference between the two disk formation scenarios discussed here. In the case where AGB wind material is captured to form a disk, one expects the abundances of the jet and AGB outflows to be the same. In the disrupted companion scenario, the disk/jet will have formed from a different star in a different evolutionary state. Thus we would expect abundance differences between the AGB and jet outflows. It may, therefore, be worthwhile for observers to search for abundance gradients between different components of multipolar flows.

We note, finally, that this paper comprises a step beyond Blackman et al. (2001b) in establishing the efficacy of MHD paradigms for PPNs/PNs in which strong magnetic fields play a role in both launching and collimating the flows. Future models should attempt to include a more detailed description of the physics of MCL disk launching in these systems, including time-dependent models (see for example von Rekowski 2004).

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