Why Magnetic Fields Cannot Be the Main Agent Shaping Planetary Nebulae

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ABSTRACT. An increasing amount of the literature reports the detection of magnetic fields in asymptotic giant branch (AGB) stars and in central stars of planetary nebulae (PNe). These detections lead to claims that the magnetic fields are the main agent shaping the PNe. In this paper, I examine the energy and angular momentum carried by magnetic fields expelled from AGB stars, as well as other physical phenomena that accompany the presence of large-scale fields, such as those claimed in the literature. I show that a single star cannot supply the energy and angular momentum if the magnetic fields have the large coherent structure required to shape the circumstellar wind. Therefore, the structure of nonspherical planetary nebulae cannot be attributed to dynamically important large-scale magnetic fields. I conclude that the observed magnetic fields around evolved stars can be understood with respect to locally enhanced magnetic loops, which can have a secondary role in the shaping of the PN. The primary role, I argue, rests with the presence of a companion.

1. INTRODUCTION

Only a small fraction of the observed planetary nebula (PN) population has spherical structure. Most PNe possess axisymmetrical structures, with a rich variety of morphologies. One of the major open questions regarding PNe is the mechanism responsible for their shaping (Balick & Frank 2002). The key question in any model is whether a binary companion, stellar or substellar, is required for the shaping of PNe, or whether stellar rotation and magnetic fields can be alone responsible for the observed morphologies.

In a series of papers (summarized in Soker 2004), I have claimed that a binary companion, stellar or substellar, must be present to form a nonspherical PN, with different binary companions influencing the mass-loss geometry from PN progenitors in different ways (e.g., accreting mass and blowing jets, spinning up the envelope, or colliding with the asymptotic giant branch (AGB) core; § 2 of Soker 2004).

On the other hand, magnetic fields are still perceived in the literature as plausible physical mechanisms that can be the sole mechanism responsible for the shaping of PNe. This idea has been reinforced by the discovery of magnetic fields in AGB stars (e.g., Etoka & Diamond 2004; Bains et al. 2004b) and more recently in actual central stars of PNe (Jordan et al. 2005).

In this paper, I explain how these findings, interesting though they are, remain far from convincing in providing an explanation of how PNe acquired their axisymmetric shapes. In particular, angular momentum and energy considerations lead to the conclusion that global magnetic fields, such as those claimed in the literature, cannot be present in AGB stars (§ 3.1). I also show that other physical considerations, such as magnetic energy (§ 3.2), magnetic stress (§ 3.3), and ambipolar diffusion (§ 3.4), make PN shaping by magnetic fields unlikely. Finally, I conclude in § 4, arguing that local magnetic fields are consistent with those observed, and that they can take a side role in the shaping of local regions inside the main nebula. The burst of strongly polarized OH maser emission in the proto-PN OH 17.7–2.0 supports a local nature of the polarized maser emission (Szymczak & Gerard 2005). The most likely primary shaping agent, however, remains the presence of companions. The binary model is supported by recent claims of more PNe harboring central binary systems (De Marco et al. 2004; Hillwig 2004; Sorensen & Pollacco 2004).

2. MAGNETIC FIELDS

2.1. The Detection of Magnetic Fields

There are now several detections of magnetic fields in PNe and around AGB stars. The most convincing measurements come from polarized maser emission (e.g., Zijlstra et al. 1989; Kemball & Diamond 1997; Szymczak et al. 1998; Szymczak & Gerard 2004, 2005; Miranda et al. 2001; Vlemmings et al. 2002, 2005; Diamond & Kemball 2003; Bains et al. 2003, 2004a, 2004b; Jordan et al. 2005). Crudely, these works find the strength of the magnetic field to be ~1 G in SiO maser clumps near the AGB stellar surface at AU, decreasing to ~10^{-3} G in H2O and OH maser clumps at r ~ 100–1000 AU. The geometry varies from large-scale ordered polarization to disordered fields, depending on the object. Large uncertainties may be involved in determining the magnetic field strength. For example, Richards et al. (2005) find discrepancies between the EVN/global VLBI and MERLIN data. It may even be that in some cases the strength of the magnetic field in maser spots is overestimated; i.e., non-Zeeman interpretation is pos-
sible in SiO masers (Wiebe & Watson 1998). Variable X-ray luminosity, as was found in Mira A (Karovska et al. 2005), can also hint at the presence of magnetic activity, although the field strength and geometry cannot be inferred.

More recently, Jordan et al. (2005) made the claim that magnetic fields were found in the actual central stars of PNe, as opposed to their precursors. They find conclusive evidence for magnetic fields in two central stars of PNe, with possible detections in an additional two PNe. From these detections, they conclude that magnetic fields have shaped the PNe.

Some of these papers jump from their magnetic field detections to the conclusion that magnetic fields shape PNe (e.g., Etoka & Diamond 2004; Bains et al. 2004b; Jordan et al. 2005), but more often than not, no justification of the mechanism is provided.

In Soker (2002), I showed that the detected magnetic fields do not necessarily play a role in globally shaping PNe. In that paper and in Soker & Kastner (2003), it is argued that the magnetic fields can only be of small coherence lengths; such small-scale fields might result from stellar magnetic spots or flares, or from jets blown by an accreting companion. These are different from the global fields needed to shape the entire PN. In the solar wind, magnetic pressure exceeds thermal pressure only in magnetic clouds (e.g., Yurchyshyn et al. 2001), which are formed by discrete mass-loss events from the Sun.

Soker & Kastner (2003) suggest that the maser spots with strong magnetic fields that are observed around some AGB stars might be similar in nature to the magnetic clouds in the solar wind, in that they represent local enhancements of the magnetic field. The fast variation in the polarization of the OH maser emission in the proto-PN OH 17.7−2.0 (Szymczak & Gerard 2005) seems to support a local explanation for a strong magnetic field. This magnetic field topology is drawn schematically in Figure 1. Locally strong magnetic fields formed by local dynamo action based on the convective envelope were also suggested to exist on the giant star Betelgeuse (Dorch 2004; Dorch & Freytag 2002). In both AGB stars and Betelgeuse, the locally strong magnetic fields are due to the strong convection in the envelope. Therefore, the magnetic field detected in maser spots might originate from one of the following sources, rather than being a large-scale field: (1) locally (but not globally) strong magnetic field spots on evolved stars, (2) a magnetically active main-sequence companion, and (3) an accretion disk around a companion, which amplifies the magnetic field of the accreted gas; the magnetic field can be transported to large distances by jets.

In this paper, I reinforce the argument presented in Soker (2002) by demonstrating in a physical way how magnetic fields of the type detected cannot play a dominate role in globally shaping PNe, although they might still play a role in local shaping (namely, on physical scales much shorter than the PN radius).

![Magnetic Field Topology](image)

**2.2. The Role of Magnetic Fields in Shaping PN Morphologies**

Models in which magnetic fields play a role in shaping the circumstellar matter belong to one of four categories:

1. Magnetic fields deflect the flow close to the stellar surface; i.e., the magnetic pressure and/or tension are already dynamically important on the AGB (or post-AGB) stellar surface (e.g., Pascoli 1997; Matt et al. 2000; Blackman et al. 2001; García-Segura et al. 2005).

2. The magnetic field is weak close to the AGB stellar surface, but it plays a dynamical role at large distances from the star (e.g., Chevalier & Luo 1994; García-Segura 1997; García-Segura et al. 1999).

3. Local magnetic fields on the AGB stellar surface are enhanced in cool spots and filaments. The dust formation rate, and hence the mass-loss rate, is enhanced above these cool spots and filaments (§ 2 in Soker 2004, and references therein). Magnetic fields never become dynamically important on a large scale. This mechanism can operate even for very slowly rotating AGB envelopes, and spin-up by companions even as small as planets is enough to account for the required rotational velocity. (This process may lead to the formation of moderately elliptical PNe; those that have small departures from sphericity but cannot account for lobes, jets, etc.)
4. Magnetic fields play a dominant role in the launching of jets from accretion disks, either around stellar companions or around the central star. The action of the jets shapes the PN.

Soker (1998) and Soker & Zoabi (2002) conclude that models 1 and 2 above, which invoke strong stellar magnetic field shaping, cannot work without a companion constantly spinning up the envelope. It is the companion that has the most influential effects on the mass-loss process, not the magnetic field per se. At most, the magnetic field plays a secondary role. Models for magnetic activity during the post-AGB phase have, in addition, the problem of explaining the bipolar outflow from systems in which the progenitor is still an AGB star. Another problem of models based on process 2 is trying to explain point-symmetric PNe, which result from precession. The previous arguments against the magnetic models 1 and 2 above are as follows: (1) in some models of type 1 above, the assumed rotation profile in the envelope is too fast or is not realistic. It is not realistic in the sense that a large angular velocity gradient is assumed to exist inside the envelope. But the same strong magnetic field inside the envelope enforces a uniform rotation in a short time (Soker & Zoabi 2002). The rotation is too fast in the sense that the required envelope rotational velocity is much higher than what a single AGB star can possess (Soker 1998). (2) There is no good explanation for the transition from spherical to nonspherical mass-loss geometry before the star leaves the AGB (Soker & Zoabi 2002). (3) Models of type 1 do not consider the role of radiation pressure on dust as the mechanism driving the high mass-loss rate at the end of the AGB.

In processes 3 and 4, the magnetic field plays an auxiliary role and is not directly responsible for the shaping of the PNe. In process 3, the mass-loss geometry is determined by radiation pressure on dust. The magnetic field that forms the cool spots only leads to a higher concentration of spots in the equatorial plane. This process is more similar to models based on nonradial pulsation of the envelope, where photospheric behavior shapes the wind, rather than models based on dynamical magnetic fields.

The jet shaping in category 4 is clearly different from the first three types of models:

1. The energy source of the magnetic field is different. In the first three categories, the main energy source is the nuclear burning in the AGB (or post-AGB) core. It drives the convective envelope that amplifies the magnetic field. In the fourth type (category), it is gravitational energy of the accreted gas that is partially converted to magnetic energy. This is a fundamental difference.

2. In the first three groups, the magnetic field is amplified by the AGB, or in some models by the post-AGB, star. In the fourth group, the magnetic field is amplified by an accretion disk around the companion or around the central post-AGB star.

3. In most models of the first three categories, the magnetic fields directly influence (shape) the nebular material as it leaves the AGB star, or at large distances. In the fourth category, that of jet shaping, the magnetic field influences only a small portion of the mass that ends up in the nebula (if it does at all). Only in some models of the first category do the magnetic fields of the post-AGB stars shape two fast jets that contain a small amount of mass.

4. As stated above, in the fourth class of models, the magnetic field is not a necessary ingredient; it is enough that two jets (or a collimated fast wind [CFW]) are launched, whether by magnetic fields, by thermal pressure, or by other mechanisms.

5. As I have shown in several papers (Harpaz & Soker 1994; Soker 1996; Soker & Zoabi 2002), in the first three classes, the AGB or post-AGB stars must be spun up by a stellar (or substellar, in the third class) companion. In the fourth class, the accretion disk is formed when the transferred mass has a high enough specific angular momentum resulting from the orbital motion.

6. As stated above, models for launching jets without magnetic fields do exist (e.g., Soker & Lasota 2004). But even if magnetic fields play a role in launching jets from accretion disks (as most researchers believe), so does the viscosity in the disk. Still, we do not term the jet shaping “viscous shaping.”

In young stellar objects, active galactic nuclei, and many other systems, it is well established that accretion disks can give rise to two jets. How exactly the two jets are launched is still an open question, but it is also beside the point made here. What matters is that the magnetic field might give rise to the jets, and that these are the jets that carve the PNe with their activity. The geometry of the jets, whether they originate at the orbiting companion or at the primary, their intensity, and their duration will be the factors determining the resulting PN shape. This process should be termed jet shaping, and not magnetic shaping.

3. PHYSICAL REASONS THAT ARGUE AGAINST A GLOBAL MAGNETIC FIELD SHAPING OF PNe

In this section, I concentrate on three key physical characteristics of AGB star winds and magnetic fields that argue against global magnetic fields playing a fundamental role in shaping PNe.

3.1. Angular Momentum Considerations

In this subsection, I demonstrate that large magnetic fields will carry angular momentum away from the stellar envelope much faster than mass is lost by the wind. This has the effect of spinning down the star on short timescales. Since the magnetic field is powered by rotation, this argument disproves the claim that such a strong large-scale magnetic field is present in the star.

At a large distance from the star, any dipole field anchored to the star (as is the Earth’s magnetic field) is negligible. The dominant component becomes the toroidal (azimuthal) component of the magnetic field carried with the wind. (The toroidal
component is the magnetic field component that is parallel to
the equatorial plane and perpendicular to both the symmetry
axis and the radial direction). If at some large radius the wind
is compressed in the radial direction, as in a shock wave, the
toroidal component, which is parallel to the shock front, in-
creases as the density (assuming the magnetic field is frozen
in to the gas). Since at a large distance the toroidal component
dominate, the magnetic pressure increases as the square of the
density. The magnetic pressure, therefore, can become larger
than the thermal pressure, as in the model of Chevalier & Luo
(1994). In their model, the wind in the transition from the AGB
to the PN phase, or the fast wind during the PN phase (here I
refer to such fields in the AGB phase as well), carries a large-
scale magnetic field. This is shown schematically in Figure 2
for one magnetic field line in the equatorial plane. Close to
the star, the magnetic pressure and tension are negligible compared with the ram pressure and thermal pressure of the wind. As
the wind hits the outer PN shell, which is the remnant of the
slow wind, it goes through a shock and slows down, and the
toroidal component of the magnetic field increases substan-
tially. This may result in the magnetic tension and pressure
becoming the dominant forces near the equatorial plane. In
particular, the magnetic tension pulls toward the center and
reduces the effective pressure in the equatorial plane. The force
due to the tension is represented by double-line arrows in Figure 2. According to this model (Chevalier & Luo 1994; García-
Segura 1997), the equatorial plane will thus be narrow, leading
to an elliptical or bipolar PN. It must be noted with regard to
this and similar models that to globally shape circumstellar
medium, the toroidal (azimuthal) field component \( B_{\theta} \) must have
the same sign around the symmetry axis.

Magnetic fields carry angular momentum; namely, the matter
attached to the field lines at large distances from the star carries
angular momentum transferred to it by the magnetic stress.
This is connected to the force that results from tension not
having a pure radial component, as seen in Figure 2. Let \( B_{\theta} \) be
the magnetic field radial component. The magnetic field carries
an angular momentum flux (angular momentum per unit area
per unit of time away from the star; e.g., Mestel 1968)

\[
\frac{dj}{dt} = \frac{\Omega B_{\theta} B_{r}}{4\pi},
\]

where \( \Omega = r \sin \theta \) is the distance from the rotation (symmetry)
axis, and \( r \) is the distance from the stellar center. I take the magnetic field components in the equatorial plane at distance \( r \) from the stellar center from Chevalier & Luo (1994):

\[
B_r(r, \theta) = B_\theta \left(\frac{R_*}{r}\right)^2 \tag{2}
\]

and

\[
B_\theta(r, \theta) = B_\theta \sin \theta \left(\frac{v_{\infty}}{v_w} \left(\frac{R_*}{r}\right)^2 \left(\frac{r}{R_*} - 1\right)\right), \tag{3}
\]

where \( v_w \) is the terminal wind speed, assumed to be in the radial direction, \( \theta \) is the altitude, measured from the polar direction (\( \theta = 0, \pi \) in the polar directions), \( B_\theta \) is the field at the stellar surface, and \( R_* \) is the stellar radius. One magnetic field line of this field in the equatorial plane is drawn in Figure 2 for \( v_{\infty} = 0.1v_w \). On the stellar surface, the field is radial. The toroidal (azimuthal, tangential) component of the magnetic field is created by the stellar rotation, which exerts tangential force on the gas and magnetic field. I consider now only the angular momentum carried by the magnetic field; the gas itself can carry additional angular momentum from the star. Substituting \( B_r \) and \( B_\theta \) in equation (1) at a large distance from the star \( (r \gg R_*) \) gives

\[
\frac{d\mathbf{J}}{dt} = \frac{B_\theta^2 R_*^2 \sin^2 \theta v_{\infty}}{4\pi v_w \left(\frac{R_*}{r}\right)^2}. \tag{4}
\]

We take the region where the magnetic field is amplified by the stellar dynamo to be within an angle \( \alpha \) from the equatorial plane. Integrating over the outgoing magnetized wind from \( \theta = \pi/2 - \alpha \) to \( \theta = \pi/2 + \alpha \) gives the total rate of angular momentum carried by the wind:

\[
\frac{d\mathbf{J}}{dt} = \int d\theta \left(\begin{array}{c}
\sin \alpha - \frac{1}{3} \sin^3 \alpha \\
\end{array}\right) B_\theta^2 R_*^3 v_{\infty} \sin \theta, \tag{5}
\]

Most of the contribution to the angular momentum comes from regions near the equatorial plane. For \( \alpha = 30^\circ \), the value of the term in parentheses is \( \sim 0.5 \), which can be considered as the scaling value.

The field measured by the observations cited above is the magnetic field at a large distance from the star \( (r \gg R_*) \), where the tangential component dominates. Let \( B_{\theta,\text{obs}} \) be the observed magnetic field at radius \( R_{\text{obs}} \), from the star, and let us use equation (3) for the tangential component, which is the component that dominates at large radii. We also note that the exact form we take for the variation of the magnetic field in the wind is not important. To globally shape the nebula, the azimuthal component must be in the same direction around the star, and this component in a dynamo model results from the stellar rotation. Therefore, eliminating \( B_\theta \) from equation (3) for \( r = R_{\text{obs}} \gg R_* \), and substituting in equation (5) for \( \theta \sim 90^\circ \), gives

\[
\frac{d\mathbf{J}}{dt} \approx 0.5B_{\theta,\text{obs}}^2 R_{\text{obs}}^2 R_* v_w \left(\frac{v_{\infty}}{v_w}\right) \tag{6}
\]

Under the assumption of a uniformly rotating envelope, which is likely to be the case in deep convective AGB stellar envelopes, the angular momentum of the stellar envelope is

\[
J_s = \eta M_{\text{env}} R_* v_{\text{rot}}, \tag{7}
\]

where \( M_{\text{env}} \) is the envelope mass, and \( \eta \) is defined such that \( \eta M_{\text{env}} R_*^2 \) is the moment of inertia of the envelope; for AGB stars, \( \eta = 0.2 \).

As explained above, I aim to demonstrate that a star with a global magnetic field will spin down on short timescales, and this will reflect negatively back onto the magnetic field, which is sustained by fast rotation. To derive a transparent expression for the spinning-down time of the envelope, I first define the mass-loss rate timescale:

\[
\tau_m = M_{\text{env}}/\dot{M}, \tag{8}
\]

where \( \dot{M} \) is the mass-loss rate from the star. This is a measure of the time the star takes to lose its envelope.

Next, I define the ratio of the magnetic energy density \( E_B = B^2/8\pi \) to the kinetic energy density \( E_k = \rho_w v_w^2/2 \), where \( \rho_w = M/4\pi v_w R_{\text{obs}}^2 \) is the wind mass density at \( R_{\text{obs}} \):

\[
\chi = \frac{E_B}{E_k} = \frac{B_{\theta,\text{obs}}^2 R_{\text{obs}}^2}{M v_w},
\]

\[
= 3.6 \left(\frac{B_{\theta,\text{obs}} R_{\text{obs}}}{1 \text{ G AU}}\right)^2 \left(\begin{array}{c}
\frac{M}{10^{-6} M_\odot} \text{ yr}^{-1}
\end{array}\right) \left(\begin{array}{c}
\frac{v_w}{10 \text{ km s}^{-1}}
\end{array}\right)^{-1}, \tag{9}
\]

where quantities are scaled with values typical for AGB stars. Another parameter is the thermal-to-magnetic pressure ratio \( \beta = E_k/E_B \). If either \( \chi \ll 1 \) or \( \beta \gg 1 \), the magnetic field is dynamically unimportant. If, on the other hand, \( \chi \) is not too small and \( \beta \) is not too large, then the magnetic field has the potential to dominate the motion of the gas.

The spinning-down time as a result of the angular momentum carried by the magnetic field is equal to the angular momentum of the envelope divided by the rate at which the magnetic field carried the angular momentum away:

\[
\tau_s = J_s \left(\frac{dJ}{dt}\right)^{-1} \approx 0.4\tau_m \chi^{-1}, \tag{10}
\]

where \( \eta = 0.2 \) was substituted, and where equations (6) and (7) were substituted in and the terms were rearranged to make the role of \( \tau_m \) and \( \chi \) explicit. The slowing time derived above
does not refer to any particular evolutionary phase. However, it is most relevant to examine the phase for which models ought explain axisymmetrical mass loss. This is when the AGB (or post-AGB) star loses its last 0.1–0.3 to explain axisymmetrical mass loss. This is when the AGB it is most relevant to examine the phase for which models ought does not refer to any particular evolutionary phase. However, the AGB slowing down, refers to the formation phase of the PN shell; it should not be compared at all with the lifetime of the descendant PN.

Here I elaborate on the meaning of equation (10). The significant quantity for the formation of the asymmetrical nebula is the fraction of the envelope mass that is lost in a nonspherical geometry. What equation (10) shows is that after losing a small fraction of the envelope, the star will spin extremely slowly, and the rest of the envelope, which contains much more mass than the mass that has been lost in a nonsymmetrical geometry, will be lost in a spherical geometry, according to the magnetic model I criticize here. To show that, I rearrange equation (10) and write it explicitly using equations (7) and (8):

\[
\frac{d}{dt} \left( \eta M_{env} R_{rot} \right) = \frac{\chi}{0.4} \frac{dM_{env}}{dt} \left( \eta M_{env} R_{rot} \right). 
\] (11)

Assuming that the envelope radius and structure, as expressed in \( \eta \), do not change much during the mass-loss process on the upper AGB, the last equation is simplified to

\[
\frac{d}{dt} \left( M_{env} v_{rot} \right) \frac{dt}{dM_{env}} = \frac{\chi}{0.4} v_{rot}. 
\] (12)

Performing the derivative gives

\[
v_{rot} + M_{env} \frac{dv_{rot}}{dM_{env}} = \frac{\chi}{0.4} v_{rot}, 
\] (13)

which, with the initial condition of rotational velocity \( v_{rot,0} \) when the envelope mass is \( M_{env,0} \), has the solution

\[
v_{rot} = v_{rot,0} \left( \frac{M_{env}}{M_{env,0}} \right)^{2\chi-1}. 
\] (14)

For typical values quoted in recent literature, \( 2.5\chi - 1 \approx 6.5 \). Substituting this value in equation (14) shows that after only 10% of the envelope has been lost, its rotation velocity drops by a factor of 2. For a loss of 30% of the envelope, the rotation velocity decreases by a factor of 10. To summarize, what equation (14) (or eq. [10]) shows is that independently of the mass-loss timescale, the envelope spins down substantially before most of the envelope mass has been lost. This implies that according to the magnetic model criticized here, most of the envelope mass will be lost in a spherical geometry.

The slowing-down timescale of the stellar rotation will be even shorter than that given by equation (10), because the wind mass itself carries angular momentum. This is true for both magnetized wind closer to the equatorial plane and for wind flowing closer to the polar direction, where the magnetic field was assumed to be weak. According to equation (3), as gas flows outward the angular momentum carried by the magnetic field increases, as the gas transfers angular momentum to the magnetic field. When nonmagnetized wind is considered, the specific angular momentum of AGB envelopes decreases as \( l \sim M_{env}^2 \) (Harpaz & Soker 1994). What equation (10) shows is that strong large-scale magnetic fields, such as those claimed in some observations and required in some models, will slow down the rotation of the AGB stellar envelope on an even shorter timescale.

Bains et al. (2003) find \( \chi \approx 3 \) in the proto-PN OH 17.7−2.0, and Bains et al. (2004a) find \( \chi \approx 4 \) in the proto-PN IRAS 20406+2953. Vlemmings et al. (2002) find the ratio of thermal to magnetic pressure to be \( \beta = 0.05 \) for a gas at a temperature of 1000 K in H$_2$O maser clumps; from that and for a wind velocity of 10 km s$^{-1}$, I find \( \chi = 5 \). In the H$_2$O maser clumps in the wind of the red supergiant S Per, Richards et al. (2005) find the thermal-to-magnetic pressure ratio to be \( \beta = 0.05 \) for a gas at a temperature of 1000 K. For a wind velocity of 10 km s$^{-1}$, I find \( \chi = 5 \). In any case, in S Per the magnitude and direction of the magnetic field measured from OH masers contains no large-scale structure and hence cannot produce large-scale shaping. It is more likely to be consistent with local magnetic clouds, as drawn in Figure 1 (Soker & Kastner 2003). Bains et al. (2003, 2004b) and other authors conclude that \( \chi > 1 \) is evidence for global shaping of the circumstellar gas by magnetic fields. However, when seen in the context of equation (10), \( \chi > 1 \) simply points to the fact that the spin-down time \( \tau_{sd} \) is smaller than the timescale over which the entire envelope is ejected. In other words, if such a field exists, the star will dramatically slow down its rotation and quench its global field before the field has time to shape the mass loss. Without a continuous supply of angular momentum, the star will not be able to sustain a strong tangential field, which is required to shape the wind in magnetic-shaping models. Such a source of angular momentum could easily be in the form of a binary companion (the small giant’s core cannot store enough angular momentum). But once we invoke the presence of a companion, other effects will come into play that will outperform the action of the magnetic field (e.g., by accreting mass and blowing jets; Soker 2004). I emphasize that the need for a companion comes from the short spinning-down time; namely, the star substantially slows down before much of the envelope is lost (\( \tau_{sd} \ll \tau_{en} \)). The spin-down results from the angular momentum carried by the gas in the wind (e.g., Harpaz & Soker 1994), or from the angular momentum due to the magnetic stress, as shown in the present paper. The inequality \( \tau_{sd} \ll \tau_{en} \), which implies the need for a companion, is better fulfilled in dynamical magnetic models when \( \chi \approx 1 \), as
evident from equation (10). Without a companion, the star will therefore spin much too slowly to expel axisymmetrical wind; a spherical rather than bipolar or elliptical PN will be formed without a companion.

In another type of model (e.g., Matt et al. 2000 for PNe; Matt & Balick 2004 for the extremely massive star η Carinae), the star has a very strong dipole magnetic field. In the model of Matt & Balick (2004), the magnetic field pressure on the photosphere is globally (not locally, as in solar spots) larger than the thermal pressure (which makes their value of the magnetic field suspiciously unrealistic). The very strong magnetic field keeps the wind in corotation with the envelope, to several stellar radii. Hence, the specific angular momentum of the wind is very high, much more than that of a wind in nonmagnetized stars. So even if such large, frozen-in dipole magnetic fields did exist, from angular momentum considerations it is clear that the star slows down on a very short timescale, such that the magnetic field could no longer be supported.

3.2. Energy Considerations

Another way to look at the requirement for a stellar companion in models based on dynamically important magnetic fields is by way of energy considerations. Away from the star, the dominant field has to be tangential, as argued in the last subsection; hence, the observed milligauss fields, which are known to originate far from the stellar surface, cannot be large-scale radial fields. As we have discussed above, the tangential field is created by rotation. The energy carried in the tangential component of the field within a wind mass element $\Delta M_w$ is

$$\Delta E_B = \frac{\chi}{2} \Delta M_w v_w^2,$$

where as before, $\chi$ is the ratio of magnetic-to-kinetic energy density. The total rotational energy of the envelope is

$$E_{\text{env}} = \frac{\eta}{2} M_{\text{env}} v_{\text{rot}}^2,$$

where a uniform envelope rotation is assumed and, as before, $v_{\text{rot}}$ is the equatorial rotation velocity. The fraction of the envelope mass that can be lost within a strong magnetic field is given by equating the energies in the last two expressions:

$$\frac{\Delta M_w}{M_{\text{env}}} = \frac{\eta}{\chi} \left( \frac{v_{\text{rot}}}{v_w} \right)^2.$$

For AGB stars, $\eta \sim 0.2$. Therefore, in order to expel at least $\sim 10\%$ of the AGB envelope mass with a strong magnetic field of $\chi \sim 1–3$, the envelope equatorial rotation velocity must be $v_{\text{rot}} \geq v_w \sim 10$ km s$^{-1}$. Such a rotation velocity is not indigenous to singly evolved AGB stars and requires the AGB to be spun up by a companion at least as massive as a brown dwarf (Harpaz & Soker 1994).

In the numerical calculations of García-Segura et al. (2005), an AGB star of radius 4.5 AU expels toroidal magnetic field, with a total ejected energy of $\sim 10^{46}$ ergs after $\sim 1000$ yr. As the global toroidal field results from rotation, the parameters of such an envelope are impossible to meet for an AGB star; using equation (16), an envelope mass of $1 M_\odot$ is required to rotate with an equatorial velocity of $70$ km s$^{-1}$—above the breakup speed. Even for a binary companion to spin up the envelope, such parameters are on the extreme.

3.3. Deceleration by Magnetic Stress

In some models of magnetic field shaping (model family 2 in § 2.1), equatorial magnetic stress slows down the wind expansion in the equatorial plane, leading to the formation of an elliptical PN (Chevalier & Luo 1994). In the Chevalier & Luo (1994) model, the stress becomes important at late times, when the post-AGB wind, which carries the magnetic field, hits the previously ejected AGB wind. This is not the case for the claimed value of the magnetic fields cited in § 2.1, where the magnetic-to-kinetic energy density is already $\chi \geq 1$ in the AGB or proto-PN phase. For a global shaping, the tangential component of the magnetic field, which is the dominant component, must encircle the star. The stress force per unit of volume in the equatorial plane, pulling toward the center, is

$$f = \frac{B^2}{4\pi R}.$$

The deceleration of the wind (or acceleration toward the center) is

$$a = f \rho = \chi \frac{v_w^2}{R}.$$

The time required to stop the wind’s expansion is

$$\tau_w = \frac{v_w}{a} = \frac{R}{\chi \frac{v_w}{R}} \approx 160 \left( \frac{R}{1000 \ \text{AU}} \right) \left( \frac{v_w}{10 \ \text{km s}^{-1}} \right)^{-1} \left( \frac{\chi}{3 \ \text{AU}} \right)^{-1} \text{yr}.$$

For the claimed value of $\chi = 3$ made by, e.g., Bains et al. (2003, 2004b; see § 3.1), the wind should have stopped long before it reached its present observed maser spots, which requires a time of $R/v_w$. In any case, this region should not expand, but rather move toward the center. This is in contradiction to the observed expansion velocity of $\sim 14$ km s$^{-1}$ found by Bains et al. (2003) in the proto-PN OH 17.7–2.0, where they also find $\chi \sim 3$. 

2006 PASP, 118:260–269
At large distances from the star, the tangential component becomes the dominant one, as is assumed throughout this paper. It is not possible that the observed magnetic field at \( r \sim 1000 \) AU and magnitude of \( \sim 10^{-3} \) G is a large-scale radial field. This is because the radial component decreases as \( r^{-2} \), which implies that the magnetic field on the photosphere at \( R_\odot \sim 1 \) AU is \( B_r \sim 1000 \) G, and the magnetic pressure \( P_B \sim 4 \times 10^4 \) ergs cm\(^{-3} \) \( \gtrsim 10 P_{\text{th},\text{s}} \), where \( P_{\text{th},\text{s}} \) is the photospheric thermal pressure, which for AGB stars is \( \sim 10^7 \) ergs cm\(^{-3} \). For small magnetic loops, on the other hand, local evolution, such as magnetic stress, can make the radial component comparable to the tangential one.

### 3.4. Ambipolar Diffusion

The magnetic field lines carry the ionized particles and diffuse through the neutral particles in a process termed ambipolar diffusion. If ambipolar diffusion is taking place, the effect of the magnetic field in directing the motion of the gas and determining the PN shape is substantially lessened. In this section, we show that AGB winds are within the ambipolar diffusion regime, and that this constitutes another reason, although not the main one, why global field shaping does not work.

Bains et al. (2004b) use the ambipolar diffusion time as given by Hartquist & Williams (1989). However, Hartquist & Williams (1989) study self-gravitating clouds supported by the magnetic field; i.e., the force due to magnetic pressure equals gravity. Therefore, Bains et al. (1989) study self-gravitating clouds supported by the magnetic field; i.e., the force due to magnetic pressure equals gravity. This is far from the situation in AGB stars, where away from the stellar surface, gravity is negligible. Therefore, Bains et al. (2004b) derive an ambipolar timescale that is several orders of magnitude too long.

The ambipolar drift velocity between ions and neutral particles is given by \( v_d = B^2/(4\pi\rho I_\rho R) \), where \( \rho_1 \) and \( \rho_\text{n} \) are the densities of the ions and neutral particles, respectively, \( R \) is the curvature radius of the magnetic field, and \( I_\rho \) is a coefficient that depends on the drift velocity itself (e.g., Shu et al. 1987). At low temperatures dominating the maser clumps, the ions frozen to the field are mostly from singly ionized metals, rather than hydrogen or helium, which have a much higher ionization potential. For the conditions relevant to the maser clumps, the drift velocity is

\[
v_d \approx 1 \left( \frac{B}{10^{-3} \text{ G}} \right)^2 \left( \frac{R}{10^{16} \text{ cm}} \right)^{-1} \left( \frac{n_1 n_\text{n}}{10^2 \text{ cm}^{-3}} \right)^{-1} \\
\times \left( \frac{\Gamma}{10^{16} \text{ cm}^3 \text{ g}^{-1} \text{ s}^{-1}} \right)^{-1} \text{ km s}^{-1},
\]

where the ion \((n_1 \sim 1 \text{ cm}^{-3})\) and neutral particle \((n_\text{n} \sim 10^4 \text{ cm}^{-3})\) number densities were scaled according to Bains et al. (2004b).

Equation (21) demonstrates that the ambipolar diffusion speed can be comparable to the wind speed for a magnetic field of \( B \sim 3 \) mG. The wind density decreases as \( r^{-2} \), and it is assumed that the same dependence holds in maser spots; hence, \( n_1 \) and \( n_\text{n} \) decrease as \( r^{-2} \). The tangential magnetic component, on the other hand, decreases as \( r^{-1} \). Therefore, the ambipolar speed increases as \( \sim r \) and cannot be neglected in models where magnetic fields are strong. In the ambipolar diffusion process, the neutral gas leaves the magnetic fields and is not influenced by it; hence, PN shaping cannot result in the ambipolar regime. I should also note in passing that the acceleration of the magnetic field lines toward the center (§ 3.3) increases.

In the local magnetic field geometry (model 3 in § 2.2), we find different parameters. Magnetic fields are expected to be weaker and their radius of curvature smaller. The ambipolar diffusion might be important in this regime (eq. [21]), and the magnetic field has enough time to further arrange itself, via stress and diffusion, into local clumps. These clumps are observed as dense maser clumps. As discussed by Soker (2002), the magnetic pressure of small-scale magnetic fields (e.g., loops) can be enhanced at large distances from the star and become greater than the thermal pressure. But it is in small, isolated regions that the magnetic fields are strong; hence, they cannot globally shape the nebula.

### 4. Discussion and Summary

In the course of the present work, I examined recent claims for strong magnetic fields around evolved stars (e.g., Miranda et al. 2001; Vlemmings et al. 2002, 2005; Diamond & Kembell 2003; Bains et al. 2003, 2004a, 2004b; Szymczak & Gerard 2004, 2005; Jordan et al. 2005) and the conclusion by some authors that these magnetic fields have had a role in shaping their PNe. By building on the arguments of Soker & Zoabi (2002) and Soker (2004), I expose additional reasons why these conclusions are unjustified.

In order to influence the large-scale structure of the circumstellar matter (the wind), the magnetic field must have a large-scale structure. However, I showed (§ 3.1) that if the observed magnetic fields have a large-scale structure, then they contain more angular momentum and energy than a single star can supply.

I have also considered deceleration by magnetic stress. Here the toroidal field around the AGB star would contribute to decelerating the wind at the equator, possibly reversing its motion. This is contrary to the observation of the wind expansion velocity. Finally, I have determined the ambipolar diffusion speed for the regimes likely to operate in AGB stars. AGB stellar winds that contain milligauss magnetic fields at \( r \sim 1000 \) AU are well within the ambipolar diffusion regimes, and the drifting of the field through the gas is significant. This diminishes the power of the magnetic field action on the gas, because it effectively decouples the gas from the field.

I also argued that any claim that the strong observed magnetic fields demonstrate that the PN was shaped by them must consider several physical characteristics of the system that can make global fields implausible PN-shaping agents. For example, as mentioned here, in some cases the surface magnetic pressure is greater than the thermal pressure on the entire sur-
face (and not locally, as in the Sun); this does not make sense, as the photosphere of the star would be unstable. As discussed in §§ 2 and 3, the models for large-scale magnetic shaping must consider the stress of the magnetic field. When this is done, some contradictions between observations and the expected behavior of the field appear.

By demonstrating that the observed magnetic fields cannot play a global role in shaping PNe, I am not denying that magnetic fields can exist in AGB stars and their progeny. These observed fields can be attributed to local ejection events, similar to the magnetic clouds in the solar wind (Soker & Kastner 2003; Fig. 1 here). Such locally enhanced magnetic fields can be formed by a dynamo mechanism rooted mainly in the vigorous convective envelope of AGB stars. For this mechanism to operate, the AGB star does not need to rotate fast; slow rotation is sufficient and can be achieved even by an orbiting planet spinning up the envelope. Local magnetic fields can influence the wind geometry by facilitating dust formation. Locally enhanced magnetic fields on the AGB stellar surface can lead to the formation of cool spots and cool filaments, as is the case in the Sun. The dust formation rate, and hence mass-loss rate, is enhanced above these cool spots and within the cool filaments (Soker & Zoabi 2002). If spots are concentrated near the equator, then the mass-loss rate is higher in that direction. This process might lead to the formation of moderately elliptical PNe (those with small departures from sphericity), but cannot account for lobes, jets, etc. Hence, these local magnetic fields can play a role in the PN morphology, but this role is secondary to that exercised by the companion of the AGB star.

Ignoring the effects of binary companions in the shaping of planetary and related nebulae will lead to erroneous conclusions. For instance, Jordan et al. (2005) argue that their discovery of magnetic field on the surface of four central stars supports the hypothesis of magnetic shaping. In the case of systems such as EGB 5 (PN G211.9+22.6), one of the four systems they discuss, the known WD companion (Karl et al. 2003) is expected to influence the mass-loss geometry more than any possible magnetic field in AGB stars.

As I have stated and demonstrated in previous papers (Soker 2004 and references therein), the effect of companions to AGB stars, even those as small as brown dwarfs and planets, can give rise to a host of physical phenomena, resulting in shaping of the AGB mass loss and subsequent wind into the observed morphologies.

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REFERENCES

Bains, I., Gledhill, T. M., Richards, A. M. S., & Yates, J. A. 2004a, in ASP Conf. Ser. 313, Asymmetrical Planetary Nebulae III, ed. M. Meixner et al. (San Francisco: ASP), 186

Bains, I., Gledhill, T. M., Yates, J. A., & Richards, A. M. S. 2003, MNRAS, 338, 287

Bains, I., Richards, A. M. S., Gledhill, T. M., & Yates, J. A. 2004b, MNRAS, 354, 529

Balick, B., & Frank, A. 2002, ARA&A, 40, 439

Blackman, E. G., Frank, A., Markiel, J. A., Thomas, J. H., & Van Horn, H. M. 2001, Nature, 409, 485

Chevalier, R. A., & Luo, D. 1994, ApJ, 421, 225

De Marco, O., Bond, H. E., Harmer, D., & Fleming, A. J. 2004, ApJ, 602, L93

Diamond, P. J., & Kemball, A. J., 2003, ApJ, 599, 1372

Dorch, S. B. F. 2004, A&A, 423, 1101

Dorch, S. B. F., & Freytag, B. 2002, in Proc. IAU Symp. 210, Modeling of Stellar Atmospheres, ed. N. E. Piskunov et al. (San Francisco: ASP), A12 (astro-ph/0208523)

Etoka, S., & Diamond, P. 2004, MNRAS, 348, 34

García-Segura, G., 1997, ApJ, 489, L189

García-Segura, G., Langer, N., Rozycka, M., & Franco, J. 1999, ApJ, 517, 767

García-Segura, G., López, J. A., & Franco, J. 2005, ApJ, 618, 919

Harpaz, A., & Soker, N. 1994, MNRAS, 270, 734

Hartquist, T. W., & Williams, D. A. 1989, MNRAS, 241, 417

Hillwig, T. 2004, in ASP Conf. Ser. 313, Asymmetrical Planetary Nebulae III: Winds, Structure, and the Thunderbird, ed. M. Meixner et al. (San Francisco: ASP), 529

Jordan, S., Werner, K., & O’Toole, S. J. 2005, A&A, 432, 273

Karlovská, M., Schlegel, E., Hack, W., Raymond, J., & Wood, B. E. 2005, ApJ, 623, L137

Kemball, A. J., & Diamond, P. J. 1997, ApJ, 481, L111

Matt, S., & Balick, B. 2004, ApJ, 615, 921

Matt, S., Balick, B., Wingele, R., & Goodson, A. 2000, ApJ, 545, 965

Meset, L. 1968, MNRAS, 138, 359

Miranda, L. F., Gomez, Y., Anglada, G., & Torrelles, J. M. 2001, Nature, 414, 284

Pascoli, G. 1997, ApJ, 489, 946

Richards, A. M. S., et al. 2005, in Proc. 7th European VLBI Network Symp., ed. R. Bachiller et al. (Madrid: Obs. Astron. Nacional), 209

Shu, F. H., Adams, F. C., & Lizano, S. 1987, ARA&A, 25, 23

Soker, N. 1996, ApJ, 469, 734

———. 1998, MNRAS, 299, 1242

———. 2002, MNRAS, 336, 826

———. 2004, in ASP Conf. Ser. 313, Asymmetrical Planetary Nebulae III: Winds, Structure, and the Thunderbird, ed. M. Meixner et al. (San Francisco: ASP), 562 (extended ver.: astro-ph/0309228)

Soker, N., & Kastner, J. H. 2003, ApJ, 592, 498

Soker, N., & Lasota, J.-P. 2004, A&A, 422, 1039

Soker, N., & Zoabi, E. 2002, MNRAS, 329, 204

Sorensen, P., & POLLACCO, D. 2004, in ASP Conf. Ser. 313, Asymmetrical Planetary Nebulae III: Winds, Structure, and the Thunderbird, ed. M. Meixner et al. (San Francisco: ASP), 515

Szymczak, M., & Gerard, E. 2004, A&A, 423, 209

———. 2005, A&A, 433, L29

Szymczak, M., Cohen, R. J., & Richards, A. M. S. 1998, MNRAS, 297, 1151
Vlemmings, W. H. T., Diamond, P. J., & van Langevelde, H. J. 2002, A&A, 394, 589
Vlemmings, W. H. T., van Langevelde, H. J., & Diamond, P. J. 2005, A&A, 434, 1029
Wiebe, D. S., & Watson, W. D. 1998, ApJ, 503, L71

Yurchyshyn, V. B., Wang, H., Goode, P. R., & Deng, Y. 2001, ApJ, 563, 381
Zijlstra, A. A., te Lintel Hekkert, P., Pottasch, S. R., Caswell, J. L., Ratag, M., & Habing, H. J. 1989, A&A, 217, 157