BACKFLOW IN POST-ASYMPTOTIC GIANT BRANCH STARS

Noam Soker
Department of Physics, University of Haifa at Oranim
Oranim, Tivon 36006, ISRAEL
soker@physics.technion.ac.il

ABSTRACT

We derive the conditions for a backflow toward the central star(s) of circumstellar material to occur during the post-asymptotic giant branch (AGB) phase. The backflowing material may be accreted by the post-AGB star and/or its companion, if such exists. Such a backflow may play a significant role in shaping the descendant planetary nebula, by, among other things, slowing down the post-AGB evolution, and by forming an accretion disk which may blow two jets. We consider three forces acting on a slowly moving mass element: the gravity of the central system, radiation pressure, and fast wind ram pressure. We find that for a significant backflow to occur, a slow dense flow should exist, such that the relation between the total mass in the slow flow, $M_i$, and the solid angle it covers $\Omega$, is given by $(M_i/\beta) > 0.1M_\odot$, where $\beta \equiv \Omega/4\pi$. The requirement for both high mass loss rate per unit solid angle and a very slow wind, such that it can be decelerated and flow back, probably requires close binary interaction.

Key words: Planetary nebulae: general – stars: AGB and post-AGB – stars: mass loss – circumstellar matter
1. INTRODUCTION

More and more supporting observations (Sahai & Trauger 1998; Kwok, Su, & Hrivnak 1998; Hrivnak, Kwok, & Su 1999; Kwok, Hrivnak, & Su 2000; Huggins et al. 2001) and theoretical considerations (Soker 1990; Soker 2001) are accumulated in support of the view that significant shaping of the circumstellar material takes place just before, after, and during the transition from the asymptotic giant branch (AGB) to the planetary nebula (PN) phase. Both the wind and radiation properties are significantly changed during these stages. As the star is about to leave the AGB the mass loss rate increases substantially, up to $\sim 10^{-5} - 10^{-4} M_\odot \text{yr}^{-1}$. This wind was termed superwind by Renzini (1981); because of confusion with other winds termed superwinds, hereafter we'll refer to this wind as FIW, for final intensive wind. After the star leaves the AGB the mass loss rate decreases down to $\sim 10^{-8} M_\odot \text{yr}^{-1}$, and its velocity increases from $\sim 10 \text{ km s}^{-1}$ to few$\times10^3 \text{ km s}^{-1}$ at the PN phase. Simultaneously, the effective temperature increases, and the post-AGB star starts to ionize the nebula around it, when by definition the PN phase starts. The changes in the mass loss rate and velocity of the wind are accompanied by a change in the wind geometry, e.g., jets and bipolar structures are formed (Soker 1990; Sahai & Trauger 1998; Kwok et al. 2000). The change in the wind geometry may result from an intrinsic processes in the AGB and post-AGB mass-losing star, or from processes in an accreting companion (e.g., Soker 2001). The increase in the wind velocity leads to collision of winds (Kwok, Purton, & Fitzgerald 1978), which can lead to instabilities during the post-AGB (proto-PN) stage (Dwarkadas & Balick 1998). As the central star starts ionizing the nebula, an ionization front propagates outward, and plays a significant role in shaping the nebula, both in the radial direction (e.g., Mellema & Frank 1995; Chevalier 1997; Schönberner & Steffen 2000), and in the transverse directions (e.g., Mellema 1995; Soker 2000b).

In the present paper we examine yet another process which may occur during the transition from the AGB to the PN stage, i.e., after the FIW ceases and before ionization starts. This is a backflow of a fraction of the dense wind toward the central star(s). This backflowing material may be accreted by the central star and/or its companion. The idea of accreting backflowing material during the post-AGB phase was raised before to explain and account for: a possible mechanism for the formation of jets from an accretion disk (Bujarrabal, Alcolea, & Neri 1998); a slower evolution along the post-AGB track (Zijlstra et al. 2000); and post-AGB stars depleted of refractory elements which compose the dust particles (e.g., Van Winckel et al. 1998) by accretion of a dust-depleted circumstellar gas (Waters, Trams & Waelkens 1992), most likely in binary systems (Van Winckel 1999).

The goal of the present paper is to explore the conditions required for a backflow to occur such that it plays a non-negligible role in the post-AGB evolution. The conditions are derived in §2, while the implications for the processes mentioned above, as well as other processes, together with a short summary, are in §3.
2. CONDITIONS FOR A BACKFLOW

In this section we derive the conditions for a backflow to occur during the post-AGB phase. We assume that a very slow flow exists along some directions, e.g., in the equatorial plane, and consider the conditions for some of this material to flow inward and be accreted by the central system, before ionization starts. After being ionized, any dense cool gas will expand and will be pushed outward by radiation and ram pressure (see below). We do not consider the deceleration of the slowly outward moving mass element, but simply assume that if there is a slowly outward moving mass element, it will reach zero radial velocity at some radius.

We therefore consider a mass-element $M_i$ with a constant density $\rho_i$ within a solid angle $\Omega$ and a radial extension $\Delta r$ at a distance $r \gg \Delta r$ from the central star, such that

$$M_i = \rho_i \Omega r^2 \Delta r. \quad (1)$$

The mass element can also be a shell where $\Omega = 4\pi$. If the sound crossing time $\Delta r/c_s$, where $c_s$ is the sound speed, is shorter than any other time scale in the process, we can take the shell to move more or less coherently. Three relevant forces are acting on the mass element in the radial direction. The gravitational force of the central star(s)

$$f_g = \frac{GM M_i}{r^2}, \quad (2)$$

where $M$ is the total mass of the central system, a binary system or a single star. The fast wind blown by the central star and its radiation push outward. The force due to the fast wind's ram pressure, assuming it is much faster than the slow wind velocity, is given by

$$f_w = \rho_w(r) v^2_w \Omega r^2 = \dot{M}_w v_w \beta, \quad (3)$$

where $\rho_w(r) = \dot{M}_w/(4\pi r^2 v_w)$, $v_w$ and $\dot{M}_w$ are the density, velocity and mass loss rate (defined positively) of the fast wind, and $\beta \equiv \Omega/4\pi$. The radiation imparts a force of

$$f_r = \frac{L_*}{c}(1 - e^{-\chi}), \quad (4)$$

where $L_*$ is the luminosity of the central system, $c$ the speed of light, and

$$\chi = \rho_i \kappa \Delta r = 7 \left( \frac{M_i}{0.01M_\odot} \right) \left( \frac{\kappa}{10 \ cm^2 \ g^{-1}} \right) \left( \frac{r}{100 \ AU} \right)^{-2} \beta^{-1}. \quad (5)$$

is the optical depth of the mass element, and we used equation (1) for the density. We scale the opacity $\kappa(r)$ with a typical value for AGB stars (Jura 1986; Winters et al. 2000), and assume that the fast wind inner to the mass element absorbs a negligible fraction of the radiation. For convenience we define three dimensionless variables. The ratio of maximum radiation pressure to the fast wind ram pressure

$$q \equiv \frac{L_*/c}{\dot{M}_w v_w} = \left( \frac{L_*/5000L_\odot}{10^{-6}M_\odot \ yr^{-1}} \right)^{-1} \left( \frac{v_w}{100 \ km \ s^{-1}} \right)^{-1}, \quad (6)$$
which does not depend on \( r \). We scale the fast wind mass loss and velocity as appropriate for a post-AGB star before it starts ionizing the nebula. The ratio of the force due to radiation and wind to that of gravity depends on \( r \), both through the dependence of optical depth \( \chi \) and gravity on \( r \).

We define \( \eta \) to be this ratio at a scaling radius \( r_0 \)

\[
\eta = \frac{\dot{M}_w v_w}{GM M_i} k(r_0) r_0^2,
\]

where

\[
k(r_0) \equiv \beta [1 + q(1 - e^{-\chi})].
\]

The equation of motion for the mass element can be written as

\[
d^2 r \over dt^2 = GM \left( \eta - {r_0^2 \over r^2} \right).
\]

We take \( r_0 \) to be the radius at which the radial velocity of the mass element is zero. When there is only gravity, the free-fall time from \( r = r_0 \) to the center \( r = 0 \), with \( v(r_0) = 0 \), is

\[
t_{ff} = \frac{\pi}{23^{1/2}} \left( \frac{r_0^{3/2}}{GM^{1/2}} \right) = 1400 \left( \frac{r_0}{400 \text{ AU}} \right)^{3/2} \left( \frac{M}{1M_\odot} \right)^{-1/2} \text{ yr}.
\]

We define the dimensionless variables

\[
\tau = t/t_{ff} \quad \text{and} \quad x = r/r_0,
\]

and write the equation of motion (9) in the form

\[
d^2 x \over d\tau^2 = \frac{\pi^2}{8} (\eta - x^{-2}).
\]

Assuming that \( \eta \) is constant and does not depend on \( x \) allows us to integrate once the last equation to give the velocity

\[
v \equiv d x \over d\tau = - \frac{\pi}{2} (\eta x + x^{-1} - 1 - \eta)^{1/2},
\]

where we substituted the initial condition \( v(1) = 0 \). This can be integrated analytically for constant values of \( \eta = 1 \) and \( \eta = 0 \). The case \( \eta = 0 \) gives the free fall solution. The time left for the object to fall from \( x \) to \( x = 0 \) is given by

\[
\tau_f(x) = \frac{2}{\pi} \left( \sin^{-1} x^{1/2} - [x(1 - x)]^{1/2} \right).
\]

As expected from our scaling \( \tau(1) = 1 \), i.e., the free fall time from \( r = r_0 \) to \( r = 0 \) is \( t_{ff} \). For \( \eta = 1 \) and with \( v(1) = 0 \), the fall time from \( x = 1 \) is infinite. This is because the wind and radiation outward-acceleration equals the gravity inward-acceleration. However, the time left to fall from \( x \) to \( x = 0 \), with \( v(1) = 0 \), is finite, and is given by

\[
\tau_f(x) = \frac{2}{\pi} \left( \ln[(1 + x^{1/2})/(1 - x^{1/2})^{-1}] - 2x^{1/2} \right).
\]
Comparing the last two equations, we find as expected, that radiation pressure and wind’s ram pressure slow down the inflow. For example, for \( \eta = 0 \), i.e., no wind and radiation pressure, the time left to fall from \( x = 0.9 \) to \( x = 0 \) is \( \tau = 0.604 \), while for \( \eta = 1 \) it is \( \tau = 1.11 \). In both cases the initial condition is \( v(1) = 0 \). These cases are less interesting, since for \( \eta = 1 \) the flow actually stagnates at \( x = 1 \). More interesting is the case of \( v(1) = 0 \) and \( 0 < \eta < 1 \), since for \( \eta > 1 \) it will be push away from the center. We numerically integrated equation (13) for these conditions and for constant values of \( \eta \). We find the backflow time to be \( \tau_f(1) = 1.28, 1.52, \) and \( 2.08 \), for \( \eta = 0.5, 0.7, \) and \( 0.9 \), respectively. For \( \eta = 0.8 \), for example, the fall back times from \( x_i = 1, 0.9 \) and \( 0.8 \) are \( \tau_f = 1.72, 1.23 \) and 0.925, respectively, where \( v(x_i) = 0 \) in these cases. For \( \eta = 0 \) the backflow times for the same initial conditions are 1, 0.85, and 0.72, respectively. The conclusion from the numerical values cited above is that the typical backflow time from \( r \sim r_0 \) is the free fall time \( \sim t_{ff} \) at \( r_0 \), but because of the radiation and wind pressures the region from which this is the backflow time is much larger than the \( \eta = 0 \) cases, extending from \( r_0 \) down to \( \sim 0.7 - 0.8r_0 \) for \( \eta > 0.8 \). So the question is what is the value of \( r_0 \) for which \( \eta = 1 \). Below and close to this radius the gas falls back in a time \( \sim t_{ff} \), while it is accelerated away for larger radii. From equation (7) we find

\[
r_0 = \left( \frac{\eta GM_i}{kM_wv_w} \right)^{1/2} = 430 \text{ AU} \left[ \frac{M}{1M_\odot} \right] \left( \frac{M_i}{0.01M_\odot} \right) \left( \frac{\dot{M}_w}{10^{-6}M_\odot} \right)^{-1} \left( \frac{v_w}{100 \text{ km s}^{-1}} \right)^{-1} \left( \frac{k}{0.1} \right)^{-1} \eta \right]^{1/2}
\]

(16)

For the backflowing mass to influence the evolution significantly, we required the fall back time to be \( \gtrsim 1000 \text{ yr} \). For the typical parameters used in equations (5) and (6) we find from equation (8) that \( k(r) \sim 2\beta \); using the time scale given by equation (10) in equation (16) gives the desired condition

\[
\frac{M_i}{\beta} \gtrsim 0.1M_\odot
\]

(17)

where as before, \( \beta \equiv \Omega/4\pi \), and \( \Omega \) is the solid angle covered by the dense backflowing material. The last condition is limited by a maximum density, since a large value of the backflowing mass \( M_i \) and small value for \( \beta \) means a very high density. We now estimate a reasonable value for the density. We assume a very slow equatorial flow, with a speed of \( v_s \sim 1 \text{ km s}^{-1} \) and with a mass loss rate per unit solid angle of \( \dot{m}_s = \dot{M}_s/4\pi \). The density of the mass elements formed by this wind is

\[
\rho_{iw} = \frac{\dot{m}_s}{r^2v_s} = 1.4 \times 10^{-15} \left( \frac{\dot{M}_s}{10^{-3}M_\odot \text{ yr}^{-1}} \right) \left( \frac{r}{400 \text{ AU}} \right)^{-2} \left( \frac{v_s}{1 \text{ km s}^{-1}} \right)^{-1} \text{ g cm}^{-3}.
\]

(18)

The minimum density is that for which the fast wind compresses the dense cool wind such that the ram pressure \( \rho_wv_w^2 \) equals the thermal pressure of the cool gas. For a molecular gas, we find this density to be

\[
\rho_{ip} = \rho_w \frac{v_w^2}{c_i^2} = 10^{-16} \left( \frac{\dot{M}_w}{10^{-6}M_\odot} \right) \left( \frac{v_w}{100 \text{ km s}^{-1}} \right) \left( \frac{r}{400 \text{ AU}} \right)^{-2} \left( \frac{\dot{T}_i}{300 \text{ K}} \right)^{-1} \text{ g cm}^{-3}.
\]

(19)

where \( c_i \) and \( T_i \) are the isothermal sound speed and temperature, respectively, of the cool gas. We find that a density of \( \rho_i \sim 10^{-15} \text{ g cm}^{-3} \) is reasonable. Using equation (1) and the definition of \( \beta \)
gives
\[
\frac{M_i}{\beta} = 0.14 \left( \frac{\rho_i}{10^{-15} \text{ g cm}^{-3}} \right) \left( \frac{r}{400 \text{ AU}} \right)^3 \left( \frac{\Delta r}{0.1r} \right) M_\odot.
\]
(20)

Condition (17) is met for the density given by equation (18), but this requires a very high mass loss rate per unit solid angle. If \( \beta = 0.1 \) this requires a total mass loss rate of \( \dot{M} = 10^{-4} M_\odot \) yr\(^{-1} \), but concentrated in particular directions, probably in the equatorial plane. All these considerations strongly suggest an equatorial dense and slow flow, such as expected in a close binary system (Mastrodemos & Morris 1999; Soker 2000a). A very fast rotation can also form such a wind (Bjorkman & Cassinelli 1993), but then a binary companion is needed to substantially spin-up the envelope.

3. IMPLICATIONS AND SUMMARY

Despite the assumptions and simplifications in deriving condition (17), we feel that the results obtained in the previous section and the implications discussed in this section are quite general. We find that for a backflow to occur on a time scale of \( t_{\text{acc}} \gtrsim 10^3 \) yr after the termination of the AGB, so that it has a non-negligible role in the post-AGB evolution, the following conditions should be met: (1) The total backflowing mass should be larger than the combined mass lost in the wind and that burned in the core. For a post-AGB mass loss rate of \( \sim 10^{-6} M_\odot \) yr\(^{-1} \) the nuclear burning is negligible, and the total required mass is \( M_{\text{acc}} \simeq 10^{-3} \left( t_{\text{acc}}/1,000 \text{ yr} \right) M_\odot \). (2) The backflowing mass should have a very low, \( \sim 1 \text{ km s}^{-1} \), terminal velocity, so that eventually it will be decelerated to zero velocity, and flow back. (3) Condition (17) on the ratio of the mass of the mass-element and the solid angle it covers \( \beta = \Omega/4\pi \) should be met. (4) For reasonable densities (eq. 18), and the required mass, we find (eq. 20) that the mass range is \( M_i \simeq 0.1 - 10^{-3} M_\odot \), and the appropriate solid angle covered by the backflowing mass is \( 1 < \beta < 10^{-2} \). We can take the typical values to be \( M_i \simeq 0.01 M_\odot \) and \( \beta = 0.1 \).

These mass loss properties required that (i) the flow be concentrated along particular directions, and (ii) have an inefficient acceleration by the stellar radiation. In a previous paper (Soker 2000a) two mechanisms which lead to such a flow were discussed. In the first mechanism proposed by Soker (2000a) magnetic cool spots (as in the Sun) are formed on the surface of slowly rotating AGB stars. The lower temperature above the spots enhances dust formation, which during the final intensive wind (FIW) may lead to an optically thick wind. This means an inefficient radiative acceleration, hence a slow flow above the spot. If the spot is small, material from the surrounding flows into the shaded region and accelerates the slow flow. However, if the spot is large, material from the surroundings will not accelerate the flow much, and it will stay slow. In the present paper we find that the condition for the slowly moving material to flow back is that the spot has \( \beta \gtrsim 0.01 \), which for a circular spot means a radius of \( R_{\text{spot}} \gtrsim 0.2 R_\ast \), where \( R_\ast \) is the stellar radius. The required mass in the slow flow is given by equation (17).
In the second mechanism a high density in the equatorial plane is formed by a binary interaction, where the secondary star is close to, but outside the AGB envelope. This mechanism is supported by observations, e.g., slowly moving equatorial gas is found around several binary post-AGB stars (Van Winckel 1999, and references therein). In the process proposed by Soker (2000a) the strong interaction between the two stars forms a dense equatorial outflow which is optically thick, leading to an inefficient radiative acceleration and a very slow equatorial flow. For a very massive and significant backflow, with a mass of $M_i \gtrsim 0.01M_\odot$ and $\beta \gtrsim 0.1$, to occur, a binary mechanism is required. For the cool spots model to form such a massive slow flow, several large spots are required. This means a strong magnetic activity, which probably requires the AGB star to be spun-up by a stellar companion. We therefore argue that in both mechanisms a binary companion is required to cause a massive flow, such that it may last for $t_{\text{acc}} \gtrsim 10^3$ yrs, possibly by as long as $\sim 10^4$ yrs in extreme cases.

Such a backflow may have the following effects. If it has a large specific angular momentum, as expected in strongly interacting binary systems, the backflowing material may form an accretion disk around one or two of the two stars. The disk(s) may blow jets or collimated fast winds (CFWs), which will play a significant role in shaping the circumstellar material. Such a possibility was briefly mentioned by Bujarrabal et al. (1998) for the proto-PN M1-92. The accreted mass may slow down the post-AGB evolution, as suggested by Zijlstra et al. (2000) for some OH/IR stars. Zijlstra et al. (2000) termed these stars retarded stars, and argue for a delay as long as $10^4$ yrs by accretion from a near-stationary reservoir. They bring supporting observations, but don’t consider the formation of such a reservoir of mass. The results of the present paper put the idea of Zijlstra et al. (2000) on a more solid ground. Another effect attributed to backflow accretion is the formation of post-AGB stars depleted of refractory elements which compose the dust particles (e.g., Van Winckel et al. 1998; Waters, Trams & Waelkens 1992; Van Winckel 1999). The separation of dust from the gas was not considered in the present paper.

Finally, we speculate on another plausible effect of the dense backflowing gas. The central star’s wind is shocked when it hits the dense material. If some dense backflowing blobs survive long into the PN phase, when the central star’s wind velocity is $\gtrsim 10^3$ km s$^{-1}$, then there will be a hard X-ray emission from the post-shock fast wind material. The dense blob will be close to the central star, making the hard X-ray emitting region hard to resolve. It is not clear if this compact hard X-ray emitting region can explain the recent Chandra observations of a “point source” in the centers of the Helix (NGC 7293) and Cat’s Eye (NGC 6543) PNe (Guerrero et al. 2000).

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