Backflow in post-asymptotic giant branch stars

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ABSTRACT

We derive the conditions for a backflow toward the central star(s) of circumstellar material to occur during the post-asymptotic giant branch (post-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 disc 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

\[
\frac{M_i}{\Omega} \approx 0.1 M_\odot, \quad \Omega = \Omega/4\pi.
\]

The requirement for both a 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, hence this process is rare.

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

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 \( \sim 10^{-5} \text{–} 10^{-4} M_\odot \text{yr}^{-1} \). This wind was termed ‘superwind’ by Renzini (1981); because of confusion with other winds termed superwinds, hereafter we will refer to this wind as ‘final intensive wind’ (FIW). 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 a few \( \times 10^3 \text{km s}^{-1} \) 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 disc (Bujarrabal, Alcolea & Neri 1998); a slower evolution along the post-AGB track (Zijlstra et al. 2001); 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 probably 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 Section 2, while the implications for the processes mentioned above, as well as other processes, together with a short summary, are in Section 3.

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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

$$M_i = \rho_i r^2 \Delta r. \tag{1}$$

The mass element can also be a shell where $\Omega = 4\pi$. If the sound crossing time $\Delta t/c_s$, 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)

$$f_g = -\frac{G M_i}{r^2}, \tag{2}$$

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, pushes outward. The force due to the ram pressure of the wind is $f_w = \rho_w v_w^2 \Omega r^2 = M_w v_w \beta$, where $\rho_w(r) = \rho_w/(4\pi r^2 v_w)$, $v_w$ and $M_w$ are the density, velocity and mass-loss rate (defined positively) of the fast wind, respectively, and $\beta = \Omega/4\pi$. The radiation imparts a force of

$$f_r = \frac{L_w}{c} \beta (1 - e^{-\chi}) \tag{4}$$

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

$$\chi = \rho_i \kappa \Delta r$$

is the optical depth of the mass element, and we used equation (1) for the density. At such a high opacity the radiation is acting on the surface of the blob facing the star. 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. We assume that in the high density gas (see below) the gas and dust flow together. Separation of dust from gas will lower the gas opacity, and will make the proposed scenario more efficient. For convenience we define three dimensionless variables. The ratio of maximum radiation pressure to the fast wind ram pressure

$$q = \frac{L_w/c}{M_w v_w} = \left( \frac{L_w}{5000 \, L_\odot} \right) \times \left( \frac{M_w}{10^{-6} \, M_\odot \, \text{yr}^{-1}} \right)^{-1} \left( \frac{v_w}{100 \, \text{km s}^{-1}} \right)^{-1}, \tag{6}$$

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 $r_0$.

$$\eta = \frac{M_w v_w}{G M_M i} k(r_0) r_0^2, \tag{7}$$

where

$$k(r_0) = \beta [1 + q(1 - e^{-\chi})]. \tag{8}$$

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

$$\frac{d^2 r}{dt^2} = \frac{GM}{r_0} \left( \frac{\eta - r_0^2}{r^2} \right). \tag{9}$$

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

$$t_{ff} = \frac{\pi}{2^{3/2} (GM)^{1/2}} = 1400 \left( \frac{r_0}{400 \, \text{au}} \right)^{3/2} \left( \frac{M}{1 \, M_\odot} \right)^{-1/2} \, \text{yr}. \tag{10}$$

We define the dimensionless variables

$$\tau = t/t_{ff} \quad \text{and} \quad x = r/r_0, \tag{11}$$

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

$$\frac{d^2 x}{dt^2} = \frac{\pi^2}{8} (\eta - x^{-2}). \tag{12}$$

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

$$v = \frac{dx}{d\tau} = -\frac{\pi}{2} (\eta x + x^{-1} - 1 - \eta)^{1/2}, \tag{13}$$

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

$$\tau(x) = \frac{2}{\pi} (\sin^{-1} x^{1/2} - [x(1-x)^{1/2}]. \tag{14}$$

As expected from our scaling $\tau(1) = 1$, 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 $x = 0.9$ to $x = 0$ is $\tau = 0.604$, while for
\[ \eta = 1 \text{ it is } \tau = 1.11. \] In both cases the initial condition is \( \zeta(1) = 0. \) These cases are less interesting, because for \( \eta = 1 \) the flow actually stagnates at \( x = 1. \) More interesting is the case of \( \zeta(1) = 0 \) and \( 0 < \eta < 1, \) because for \( \eta > 1 \) it will be pushed away from the centre. We numerically integrated equation (13) for these conditions and for constant values of \( \eta. \) We find the backflow time to be \( \tau(1) = 1.28, 1.52 \) and 2.08, for \( \eta = 0.5, 0.7 \) and 0.9, respectively. For \( \eta = 0.8, \) for the fall-back mass from \( x_1 = 1, 0.9 \) and 0.8 are \( \tau(1) = 1.72, 1.23 \) and 0.925, respectively, where \( \zeta(x_1) = 0 \) 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 \( r = r_0 \) is the free-fall time \( \sim t_0 \) 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 \( \sim 0.7-0.8 r_0 \) for \( \eta \geq 0.8. \) 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 \( \sim t_0, \) while it is accelerated away for larger radii. From equation (7) we find

\[
\frac{r_0}{e} = \left( \frac{\eta G M_{\odot}}{k M_w v_w} \right)^{1/2} = 430 \text{ au} \left( \frac{M}{1 \text{ M}_\odot} \right) \left( \frac{M_1}{0.01 \text{ M}_\odot} \right) \left( \frac{M_w}{10^{-6} \text{ M}_\odot} \right)^{-1}
\]

\[
\times \left( \frac{v_w}{100 \text{ km s}^{-1}} \right)^{-1} \frac{\kappa}{0.1}^{-1} \xi^{1/2}
\]

(16)

Note again the meaning of \( r_0. \) We assume that many dense blobs are expelled from the central star(s) at different velocities (we argue that this process requires strong binary interaction). Slow blobs will be decelerated by gravity, and stagnate at radius \( r. \) Blobs for which \( r < r_0 \) will fall back to the central star(s). The value of \( r \sim \text{ several} \times 100 \) au is much larger than the required binary separation, but we note that this is of the order of the observed size of the slowly moving gas in the Red Rectangle (e.g. Jura & Kahane 1999).

For the backflowing mass to influence the evolution significantly, we required the fall back time to be \( \sim 1000 \) 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

\[
M_1 \geq 0.1 \text{ M}_\odot
\]

(17)

where as before, \( \beta = \Omega/4\pi, \) 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 \( M_1 \) 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 a mass-loss rate per unit solid angle of \( M_1 = \beta M_1/4\pi. \) The density of the mass elements formed by this wind is

\[
\rho_w = M_1 \frac{m_1}{r^2 v_s} \sim 1.4 \times 10^{-15} \left( \frac{M_1}{10^{-3} \text{ M}_\odot \text{ yr}^{-1}} \right)
\]

\[
\times \left( \frac{r}{400 \text{ au}} \right)^{-2} \left( \frac{v_s}{1 \text{ km s}^{-1}} \right)^{-1} \text{ g cm}^{-3}.
\]

(18)

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

\[
\rho_w = \rho_s \frac{v_w^2}{c_s^2} = 10^{-16} \left( \frac{M_w}{10^{-6} \text{ M}_\odot} \right) \left( \frac{v_w}{100 \text{ km s}^{-1}} \right)
\]

\[
\times \left( \frac{r}{400 \text{ au}} \right)^{-2} \left( \frac{T_i}{300 \text{ K}} \right)^{-1} \text{ g cm}^{-3}.
\]

(19)

where \( c_s \) and \( T_i \) are the isothermal sound speed and temperature, respectively, of the cool gas. We find that a density of \( \rho_w \sim 10^{-15} \) g cm\(^{-3}\) is reasonable. Using equation (1) and the definition of \( \beta \) gives

\[
\frac{M_1}{\beta} = 0.14 \left( \frac{\rho_w}{10^{-15} \text{ g cm}^{-3}} \right) \left( \frac{r}{400 \text{ au}} \right)^{3} \left( \frac{\Delta r}{0.1r} \right) \text{ M}_\odot.
\]

(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 \( \dot{M} = 10^{-4} \text{ M}_\odot \text{ yr}^{-1}, \) but concentrated in particular directions, probably in the equatorial plane. The required mass-loss process is far from being spherically symmetric. All these considerations strongly suggest an equatorial dense, slow flow, such as is 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 in that case a binary companion is needed to substantially spin-up the envelope. We did not consider the centrifugal force, which is negligible at a distance of \( r \geq 100 \) au – about two orders of magnitude larger than the required binary separation. Even if the specific angular momentum per unit mass of the expelled gas is equal to that of the binary system, the centrifugal force is much smaller than gravity at the relevant distances. The centrifugal force becomes important only when the backflowing gas is close to the binary system, where it can form an accretion disc.

3 IMPLICATIONS AND SUMMARY

Despite the assumptions and simplifications in deriving condition (17), we feel that the results obtained in the previous section and discussed in this section are quite general. We find that for a backflow to occur on a time-scale of \( t_{\text{acc}} \simeq 10^7 \) 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:

(i) The total backflowing mass should be larger than the mass lost in the wind and that burned in the core combined. For a post-AGB mass-loss rate of \( \sim 10^{-6} \text{ M}_\odot \text{ yr}^{-1} \) the nuclear burning is negligible, and the total required mass is \( M_{\text{acc}} = 10^{-3} (t_{\text{acc}}/1000 \text{ yr}) \text{ M}_\odot \).

(ii) 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.

(iii) Condition (17) on the ratio of the mass of the mass-element and the solid angle it covers is \( \beta = \Omega/4\pi \) should be met.

(iv) For reasonable densities (equation 18), and the required mass, we find (equation 20) that the mass range is \( M_1 = 0.1-10^{-3} \text{ M}_\odot, \) and the appropriate solid angle covered by the backflowing mass is \( 10^{-2} \leq \beta < 1. \) We can take the typical values to be \( M_1 = 0.01 \text{ M}_\odot \) and \( \beta = 0.1. \)

These mass-loss properties required that (1) the dense flow be concentrated along particular directions, and (2) 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 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 surroundings of the wind above the spot (shaded region) flows into the shaded region and accelerates the slow flow. However, if the spot is large, material from the surroundings of the shaded region 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 $\beta \gtrsim 0.01$ (more than 1 per cent of the stellar surface is covered by the spot), 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, 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 $M_1 \gtrsim 0.01 M_\odot$ 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 $t_{\text{acc}} \gtrsim 10^3 \text{ yr}$, possibly by as much as $\sim 10^4 \text{ yr}$ in extreme cases. The constraints on the binary interaction, such that it causes the outflow to be far from being spherically symmetric, imply that this process is rare.

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 disc around one or two of the two stars. The disc(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, Alcolea & Neri (1998) for the proto-PN M1-92. The accreted mass may slow down the post-AGB evolution, as suggested by Zijlstra et al. (2001) for some OH/IR stars. Zijlstra et al. (2001) termed these stars ‘retarded stars’, and argued for a delay of as long as $10^4 \text{ yr}$ by accretion from a non-stationary reservoir. They bring supporting observations, but do not consider the formation of such a reservoir of mass. The results of the present paper put the idea of Zijlstra et al. (2001) on to 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. Waters et al. 1992; Van Winckel et al. 1998; Van Winckel 1999). The separation of dust from the gas was not considered in the present paper. We assumed that in the high-density flow, dust and gas flow together. Any separation of gas and dust will lower the gas opacity, and will make the proposed process more efficient.

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 wind velocity of the central star is $\gtrsim 10^3 \text{ km s}^{-1}$, 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 centres of the Helix (NGC 7293) and Cat’s Eye (NGC 6543) PNe (Guerrero et al. 2001).

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