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Wind-driven evolution of accreting binaries and the progenitors of SNe Ia

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Abstract. Strong self-excited winds from the companion are very likely in many accreting binaries, and we investigate their evolutionary effects. We show that in white dwarf systems with large mass ratios $M_2/M_1 \gtrsim 1.3$, self-excited winds naturally stabilize mass transfer at a threshold value $\dot M_{\rm b} \sim 10^{-8} M_\odot$ yr^{-1}. Near this threshold value, irradiation-induced wind loss rates from the companion star become much higher, because of intermittent stable hydrogen shell burning on the accreting white dwarf. The white dwarf can accrete a substantial mass ($\sim 0.5 M_\odot$) during this evolution, and may produce a Type Ia supernova directly in some cases. In other cases the evolution may produce supersoft X-ray binaries with quite massive white dwarfs and moderate-mass companions, which then undergo thermal-timescale mass transfer and may produce SNe Ia themselves.

We show that the stability of such 'bootstrap' evolution via a self-excited wind is indeterminate if the wind losses depend linearly on the mass transfer rate, as assumed in some analyses of low-mass X-ray binary evolution. Such analyses must be repeated with a more precise form of the wind-driving law. In practice systems with small mass ratios will tend to stabilize in regimes where the wind-driving law is slower than linear, with the opposite result for large mass ratios.

Key words: accretion, accretion disks – stars: mass-loss – novae, cataclysmic variables – supernovae: general – white dwarfs – X-rays: stars

1. Introduction

In the accompanying paper (Van Teeseling & King 1998; hereafter Paper I) we showed that the rather soft spectrum of an accreting white dwarf is very likely to drive a wind from the companion star. We showed that if the white dwarf accretes rapidly enough that the accreted hydrogen is burnt without significant expansion of the white dwarf envelope, the resulting wind losses are strong enough to drive Roche lobe overflow at a sufficiently high rate that the process is self-sustaining. Such systems will appear as supersoft X-ray binaries, at least episodically; the SMC source 1E 0035.4-7230 may be an example. In binaries with small mass ratios $M_2/M_1 \lesssim 1$ (as appropriate for 1E 0035.4-7230, whose period is only 4.1 hr) some 'accidental' event such as a long shell flash either on the cooling white dwarf, or because of accretion, is required to trigger the process. Without such an event, these systems will appear as cataclysmic variables, in which the accretion rates are well below the critical value

$$\dot M_{\rm b} = 0.63 \times 10^{-7} 10^{-0.8 m_1}\ M_\odot {\rm yr}^{-1}, \qquad (1)$$

(with m_1 the primary mass M_1 in $M_\odot$). At accretion rates below $\dot M_{\rm b}$, the white dwarf envelope expands to a radius $\gtrsim$ R$_\odot$ during a hydrogen shell flash, and may eject a significant fraction of the accreted matter (e.g. Fujimoto 1982). As a consequence, $\dot M_{\rm b}$ is roughly the minimum accretion rate for effective irradiation of the companion during the hydrogen burning phase.

In this paper we consider the general case of self-excited wind-driven evolution of accreting binaries, without restrictions on the mass ratio. We begin by performing a stability analysis, incidentally showing that the case considered in Paper I is indeed stable, as assumed without proof there. We shall also show that self-excited winds can actually stabilize mass transfer at a value $\sim \dot M_{\rm b}$ in white dwarf systems with large mass ratios $M_2/M_1 \gtrsim 1$. Some of these systems may be able to drive the white dwarf mass over the Chandrasekhar limit and produce a Type Ia supernova.

2. Stability of wind-driven mass transfer

As in Paper I we consider mass transfer at a rate $\dot M_{\rm tr} < 0$ from the secondary (mass M_2) to the primary (mass M_1). The secondary is assumed also to lose mass in a wind at a rate $\dot M_{\rm w}$, the wind mass carrying specific angular momentum β_2 times that of the secondary. For generality we will also allow for systemic angular momentum loss (e.g. via gravitational radiation) at a rate $\dot J = -J/t_J$. Then the rate of change of the secondary's Roche lobe $R_{\rm L}$ can be written

$$\frac{\dot R_{\rm L}}{R_{\rm L}} = -\frac{2\dot M_{\rm tr}}{M_2}\left[c_1 - c_2\frac{M_2}{M_1}\right] - \frac{2}{t_J} + \frac{\dot M_{\rm w}}{M M_2}\left[\left(2\beta_2 - c_3\right)M_1 - c_4 M_2\right] \qquad (2)$$

with for $M_2 \lesssim M_1$,

$$c_1 = \frac{5}{6}, c_2 = 1, c_3 = \frac{5}{3}, c_4 = 1 \tag{3}$$

(see Paper I). This relation assumes Paczyński's (1971) approximation for the secondary's Roche lobe, i.e.

$$\frac{R_{\rm L}}{a} = 0.462\left(\frac{M_2}{M}\right)^{1/3} \tag{4}$$

where a is the binary separation and $M = M_1 + M_2$, which holds for $M_2 \lesssim M_1$. For the case $M_2 \gtrsim M_1$ we use the fact that the ratio of the Roche lobe radii for M_1, M_2 can be approximated to better than 5% accuracy by

$$\frac{R_{\rm L2}}{R_{\rm L1}} = \left(\frac{M_2}{M_1}\right)^{0.45} \tag{5}$$

(King et al. 1997b), and now use the relation corresponding to Eq. (4) to express $R_{\rm L1}$ in terms of M_1, M and a. This gives the Roche lobe of M_2 as

$$\frac{R_{\rm L}}{a} = 0.462 M_1^{-0.12} M_2^{0.45} M^{-0.33} \tag{6}$$

for the case $M_2 \gtrsim M_1$. Then the constants in Eq. (2) become

$$c_1 = 0.78, c_2 = 1.06, c_3 = 1.55, c_4 = 0.88 \tag{7}$$

in place of (3). The secondary's radius R_2 behaves as

$$\frac{\dot R_2}{R_2} = \zeta\frac{\dot M_2}{M_2} + \frac{1}{t_{\rm nuc}} = \zeta\frac{\dot M_{\rm tr}}{M_2} + \zeta\frac{\dot M_{\rm w2}}{M_2} + \frac{1}{t_{\rm nuc}}, \tag{8}$$

where ζ is the effective mass-radius index, i.e. that along its actual evolutionary track, and we have allowed for possible nuclear expansion of the secondary on a timescale $t_{\rm nuc}$.

For stationary mass transfer we require $\dot R_{\rm L} = \dot R_2$. If this equality does not hold, the transfer rate will evolve according to

$$\dot F = F\frac{R_2}{H}\left[\frac{\dot R_2}{R_2} - \frac{\dot R_{\rm L}}{R_{\rm L}}\right], \tag{9}$$

where $H \simeq 10^{-4}R_2$ is the scaleheight near the inner Lagrange point, and we have written $F = -\dot M_{\rm tr} > 0$ to avoid confusion over signs. Using Eqs. (2) and (8) in Eq. (9) we get

$$\dot F = aF\left[-DF - B\dot M_{\rm w} + \frac{M_2}{2t_{\rm dr}}\right], \tag{10}$$

where

$$a = \frac{2R_2}{HM_2}, \tag{11}$$

$$\frac{1}{t_{\rm dr}} = \frac{1}{t_{\rm nuc}} + \frac{2}{t_J}, \tag{12}$$

and

$$D = \frac{\zeta}{2} + c_1 - c_2\frac{M_2}{M_1}, \tag{13}$$

$$B = \left(\beta_2 - \frac{c_3}{2} - \frac{\zeta}{2}\right)\frac{M_1}{M} - \left(\frac{c_4}{2} + \frac{\zeta}{2}\right)\frac{M_2}{M}. \tag{14}$$

Eq. (10) shows that unless the square bracket S vanishes, F changes on a timescale $\sim 1/aS$. For $F \ll M_2/(2t_{\rm dr})$ this means that F rises exponentially on a timescale $\sim (H/R_2)t_{\rm dr}$. For larger F the timescale is $\sim (H/R_2)(M_2/F)$, i.e. about 10^{-4} times the current mass-transfer timescale (see the discussion in Sect. 3. below). Accordingly we can always regard M_2 and thus D, B as constant during such changes, and we shall further assume that this is true of $t_{\rm dr}$ also (this is valid for typical angular momentum loss mechanisms such as gravitational radiation and magnetic braking, and also for nuclear burning). Then Eq. (10) is a first-order ordinary differential equation for $F(t)$, and it is easy to check the stability of a given stationary solution $F = F_0$. From Eq. (10) we have

$$F_0 = -\frac{B}{D}\dot M_{\rm w}(F_0) + \frac{M_2}{2Dt_{\rm dr}}. \tag{15}$$

We consider a small perturbation $F = F_0 + u, |u/F| \ll 1$. Expanding to first order in u, Eq. (10) implies

$$\dot u = -aF_0 u\left[D + B\frac{{\rm d}\dot M_{\rm w}}{{\rm d}F}\right]. \tag{16}$$

Hence stability requires

$$D + B\frac{{\rm d}\dot M_{\rm w}}{{\rm d}F} > 0, \tag{17}$$

where all quantities are evaluated at the stationary solution $F = F_0$. Now using Eq. (15) to eliminate B gives

$$D\left[1 - \frac{{\rm d}\ln|\dot M_{\rm w}|}{{\rm d}\ln F}\right] + \frac{M_2}{2t_{\rm dr}F_0}\frac{{\rm d}\ln|\dot M_{\rm w}|}{{\rm d}\ln F} > 0 \tag{18}$$

or, again using Eq. (15),

$$D\left[1 - \frac{B\dot M_{\rm w}}{B\dot M_{\rm w} - M_2/2t_{\rm dr}}\frac{{\rm d}\ln|\dot M_{\rm w}|}{{\rm d}\ln F}\right] > 0. \tag{19}$$

In cases where the wind losses are negligible, i.e. $|B\dot M_{\rm w}| \ll M_2/2t_{\rm dr}$, this reduces to the familiar requirement

$$D > 0. \tag{20}$$

In wind-driven cases, with $|B\dot M_{\rm w}| \gg M_2/2t_{\rm dr}$, we get instead

$$D\left[1 - \frac{{\rm d}\ln|\dot M_{\rm w}|}{{\rm d}\ln F}\right] > 0. \tag{21}$$

This shows that, remarkably, wind-driven mass transfer in systems with large mass ratios ($D < 0$) is stable provided that the wind mass loss rate increases more rapidly than F itself as F increases. This stable regime is therefore only open to systems whose evolution is driven by a wind from the companion; moreover the wind must be self-excited rather than intrinsic, as the latter is by hypothesis insensitive to F. The physical reason for the stability is clear: although mass transfer tries to shrink the system ($D < 0$), the resulting increase in F tends to drive a much stronger wind from the companion, thus stabilizing the system by expanding it again. Note that the indeterminate wind-driven case $\mathrm{d}\ln|\dot{M}_\mathrm{w}|/\mathrm{d}\ln F = 1$ implies

$$\dot{M}_\mathrm{w} = -cF,\ c > 0, \tag{22}$$

so substituting in Eq. (17) gives the stability condition as

$$D - cB > 0. \tag{23}$$

In practice however no wind driving law can be precisely linear as in (22), and the stability will be determined by the deviations from the linear dependence, using the general criterion (21). We note that in considering possible 'bootstrap' evolution of low–mass X–ray binaries, Ruderman et al (1989) actually do assume a precisely linear wind law (their Eq. 3.30). They effectively use the stability criterion (23), regarding a bootstrap as established if $cB/D > 1$, which is equivalent to Eq. (23) under their assumptions $D > 0, cB > 0$. Thus mass transfer is formally unstable, and will rise to a value where the wind driving law is slower than linear. Ruderman et al implicitly assume that this first occurs near the Eddington limit. This analysis is however indeterminate because any deviation of the wind driving law from the precisely linear form assumed will allow stabilization well before this value is reached. Clearly systems with small mass ratios ($D > 0$) will tend to stabilize in regimes where the wind losses increase more slowly than linearly with F, with the opposite result for large mass ratios ($D < 0$).

We can now deal specifically with the case of the irradiation-driven winds in white dwarf binaries. From Paper I we know that these obey

$$\dot{m}_\mathrm{w} = 3.5 m_2^{5/6} m^{-1/3} (\eta_s \eta_a \dot{m})^{1/2} \phi, \tag{24}$$

for $M_2 \lesssim M_1$; using Eq. (6) we get

$$\dot{m}_\mathrm{w} = 3.5 m_2^{0.95} m^{-1/3} m_1^{-0.12} (\eta_s \eta_a \dot{m})^{1/2} \phi, \tag{25}$$

for $M_2 \gtrsim M_1$. Here $\dot{m}, \dot{m}_\mathrm{w}$ are $-\dot{M}_\mathrm{tr}$ ($= F$) and $-\dot{M}_\mathrm{w}$ in units of $10^{-7} M_\odot$ yr^{-1}, m_2, m are M_2, M in $M_\odot$; η_s measures the efficiency of the primary's spectrum in producing ionizing photons, normalized to the case of supersoft X-ray temperatures $\sim$ few times 10^5 K, and ϕ is an efficiency factor parameterizing the fraction of the companion's face which is irradiated, and the fraction of the wind mass escaping the system. Most importantly, η_a measures the luminosity per gram of matter accreted relative to the value (4.6×10^{18} erg g^{-1}, Iben 1982) for hydrogen shell burning.

For $\dot{M} \gtrsim \dot{M}_\mathrm{b}$, the accreted hydrogen is burnt in recurrent flashes, while the white dwarf envelope remains relatively small ($\lesssim R_\odot$). During these intermittent flashes, irradiation of the companion and driving of a wind from the heated side are strong. For $\dot{M} \lesssim \dot{M}_\mathrm{b}$ the flashes are more explosive, with larger expansions and possibly significant mass ejection (e.g. Fujimoto 1982). Hence for these lower accretion rates irradiation is much less efficient in driving a wind during burning flashes.

Below $\dot{M}_\mathrm{b}$, η_a is given by gravitational energy release alone as $GM_1/4.6\times10^{18}R_1$ ($R_1 =$ primary radius), while above $\dot{M}_\mathrm{b}$, η_a is unity. For $0.6 < m_1 < 1.2$, a good approximation to the white dwarf mass-radius relation (Nauenberg 1972) is

$$R_1 = 5.6 \times 10^8 m_1^{-1}\ \mathrm{cm}. \tag{26}$$

Using this we can therefore write

$$\dot{m}_\mathrm{w} = A_l \dot{m}^{1/2},\ \ \dot{m} < \dot{m}_\mathrm{b}, \tag{27}$$

$$\dot{m}_\mathrm{w} = A_h \dot{m}^{1/2},\ \ \dot{m} > \dot{m}_\mathrm{b}, \tag{28}$$

where

$$A_l = 0.8 m_2^{5/6} m_1 m^{-1/3} \eta_s^{1/2} \phi \tag{29}$$

$$A_h = 3.5 m_2^{5/6} m^{-1/3} \eta_s^{1/2} \phi \tag{30}$$

for $M_2 \lesssim M_1$, and

$$A_l = 0.8 m_2^{0.95} m_1^{0.88} m^{-1/3} \eta_s^{1/2} \phi \tag{31}$$

$$A_h = 3.5 m_2^{0.95} m_1^{-0.12} m^{-1/3} \eta_s^{1/2} \phi \tag{32}$$

for $M_2 \gtrsim M_1$. This form is shown in Fig. 1. Clearly we have $\mathrm{d}\ln|\dot{M}_\mathrm{w}|/\mathrm{d}\ln F > 1$ in the the regime near $\dot{m} = \dot{m}_\mathrm{b}$, so solutions with large mass ratio ($D < 0$) are stable here. Since for a wind-driven solution Eq. (15) implies

$$\dot{m}_\mathrm{w} = \frac{D}{B}\dot{m}, \tag{33}$$

(the ratio B/D was called g in Paper I) we see that stable solutions exist if and only if

$$\frac{\dot{m}_\mathrm{b}^{1/2}}{A_h} < \frac{B}{D} < \frac{\dot{m}_\mathrm{b}^{1/2}}{A_l}. \tag{34}$$

If this condition holds there are also two unstable solutions on each side of the stable one with $F \simeq \dot{M}_\mathrm{b}$ (see Fig. 1). If it fails, only one of these unstable solutions is available.

3. Stable solutions for accreting white dwarfs

We can now characterize the possible solutions for mass transfer allowing for self-excited wind losses in terms of the effective mass-radius index ζ and mass ratio $q = M_2/M_1$. Fig. 2 shows the ζ, q plane for positive q and $\zeta > -1/3$, the value for rapid mass loss from a fully convective (or degenerate) star. (Note that $\zeta \gg 1$ for rapid mass loss from a largely radiative star, and $\zeta \simeq 1, 0.6$ for stars near the lower and upper main sequences losing mass on timescales longer than their Kelvin times.) The curves $D = 0, B = 0$ divide this plane into 4 regions.

In the region $D < 0, B < 0$ we have possible stable solutions with $F \simeq \dot{M}_\mathrm{b}$ as we have seen in Sect. 2. If Eq. (34) holds, the system will try to find this solution, as F will increase monotonically towards the value $\dot{M}_\mathrm{b}$ once the system comes into contact. As we have seen above, the initial rise is exponential, on

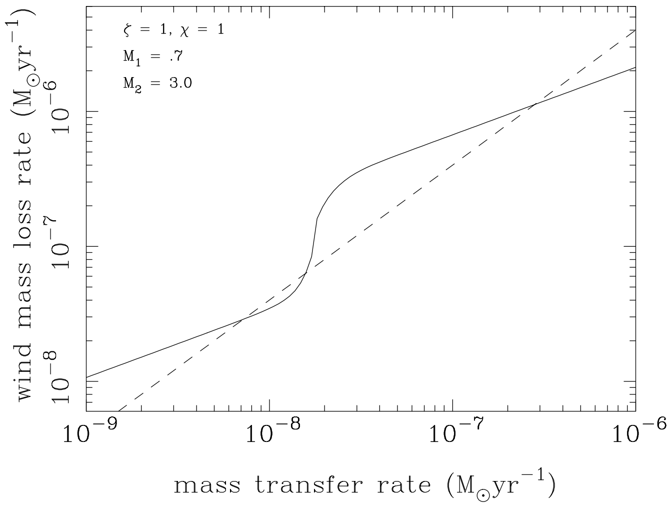

Fig. 1. The solid curve shows the wind mass-loss rate from the irradiated secondary as a function of the mass transfer rate. The dashed line shows the mass transfer rate driven by mass and angular momentum loss in the wind. A stationary solution corresponds to the intersections of the two curves. Only the central intersection on the steep part of the $\dot{m}_{\rm w} - \dot{m}$ curve is stable for large mass ratios

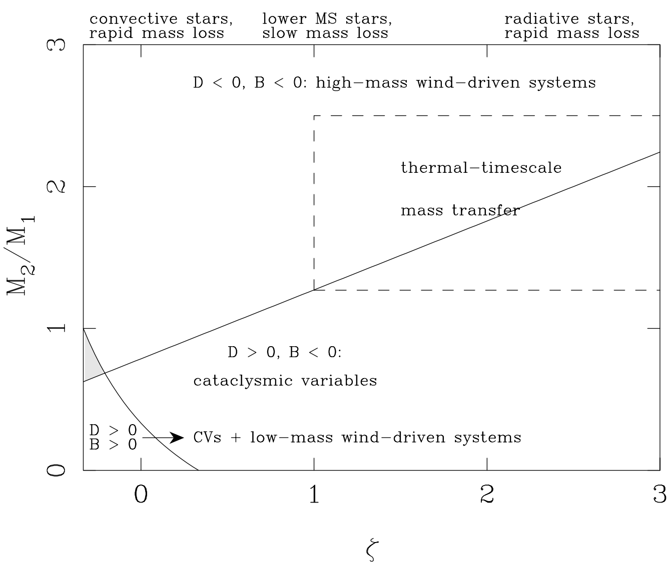

Fig. 2. The possible stable mass transfer solutions for white dwarf binaries in the $\zeta - q$ plane. No stable soutions are possible in the shaded area corresponding to $D < 0, B > 0$

a timescale $(H/R_2)t_{\rm dr} \sim 10^{-4}t_{\rm dr}$. Once F reaches values of order $M_2/(2Dt_{\rm dr})$, wind–driving will become appreciable, and F now rises as $\sim (t_0 - t)^{-1}$, with $t_0 \sim (H/R_2)(M_2/2|D|F)$. For $F < \dot{M}_{\rm b}$, this gives a timescale

$$t_0 \sim 10^4 m_1 (F/\dot{M}_{\rm b})^{-1} \text{ yr} \tag{35}$$

for the increase of mass transfer (using $|D| \sim M_2/M_1$). For comparison, the time between nova outbursts is

$$t_{\rm rec} \sim 10^4 R_9^{2.8} m_1^{-0.7} (F/\dot{M}_{\rm b})^{-1} \text{ yr} \tag{36}$$

using the scalings discussed on p. 287 of Warner (1995) and references therein: R_9 is the white dwarf radius in units of 10^9 cm. Hence the ratio of nova recurrence time to the timescale for mass transfer increase is

$$\frac{t_{\rm rec}}{t_0} \sim R_9^{2.8} m_1^{-1.7} \sim 0.2 m_1^{-4.5} \tag{37}$$

where we have used (26) at the last step. We therefore expect that flashes will recur several times before the mass transfer rate increases significantly for any white dwarf mass $M_1 \gtrsim 0.7M_\odot$. Once F reaches $\dot{M}_{\rm b}$ such systems will thus produce a continuous series of shell flashes which are stabilizing, i.e. have sufficiently small envelope expansion that irradiation is just able to drive a strong enough wind to prevent further contraction of the system under mass transfer. Systems with $D < 0, B < 0$ and $M_1 \gtrsim 0.7M_\odot$ therefore settle stably at $F \simeq \dot{M}_{\rm b}$. By contrast, in systems with low–mass white dwarfs the recurrence time for the flashes is so long that F may have increased significantly beyond the stable regime near $\dot{M}_{\rm b}$ before the next flash occurs. Such systems would thus reach the supersoft regime with $F > \dot{M}_{\rm b}$.

If $D < 0, B < 0$ but the system does not settle in the stable regime with $F \simeq \dot{M}_{\rm b}$, (e.g. Eq. (34) fails, or $M_1 \lesssim 0.7M_\odot$), mass transfer is formally unstable. However, systems with mass ratios q such that mass transfer is stable on a dynamical timescale (i.e. $D > 0$ for some $\zeta > 1$, assuming that the secondary is largely radiative) but unstable on a thermal timescale (i.e. $D < 0$ for $\zeta \simeq 1$) can transfer mass steadily at a rate set by the secondary's thermal timescale (Van den Heuvel et al. 1992). These solutions are thought to characterize many of the known supersoft X-ray binaries. In practice the defining conditions restrict q to a fairly narrow range straddling the line $D = 0$ for a range $\zeta \gtrsim 1$, shown schematically in Fig. 2. In these solutions the wind term $B\dot{M}_{\rm w}$ is of a similar order to $\dot{M}_{\rm tr}$ in modulus, and slightly reinforces the thermal-timescale mass transfer (see Paper I). If the system does not have q, ζ in this region, and is otherwise unstable, a common envelope must ensue.

The region $D < 0, B > 0$ is completely unstable, and must lead to common-envelope evolution, as all 3 terms on the rhs of Eq. (10) are positive and thus no steady state is possible. Note that since $\zeta \lesssim -0.22$ in this region, thermal-timescale mass transfer is unlikely here.

The region $D > 0, B < 0$ has stable solutions with $F \leq M_2/2Dt_{\rm dr}$, i.e. possibly reduced below the usual rate driven by angular momentum losses or nuclear expansion because of the stabilizing effect of the wind.

Finally the region $D > 0, B > 0$ has stable solutions with $F = M_2/2Dt_{\rm dr} + B\dot{M}_{\rm w}/D$. These however fall into two groups: in the first, the increase over the usual angular-momentum loss/nuclear expansion solution is very slight, i.e. $F \simeq M_2/2Dt_{\rm dr}$, because the accretion efficiency is too low to drive a strong wind (the weak wind which is present raises the mass transfer very slightly). In the second type of solution the accretion efficiency is high enough to drive a strong wind, which dominates the mass transfer, i.e. $F \simeq B\dot{M}_{\rm w}/D$. This second type of solution is the one discussed in Paper I in connection with supersoft X–ray binaries with low–mass secondaries, and evidently requires some external event to trigger it. (Tout & Hall 1991 studied this region for the case of an intrinsic wind from the secondary: note that the condition $B > 0$ implies $\zeta < 1/3$, in agreement with their result.) Hence with this exception, solutions with $D > 0$ correspond to standard cataclysmic-variable evolution, slightly modified by the weak wind from the companion.

Of course, as the system evolves, its mass ratio drops, so it may migrate across the boundaries of the various regions in Fig. 2. Similarly we shall see that in the course of evolution, the condition Eq. (34) for stable solutions in the region $D < 0, B < 0$ may fail in the sense of allowing only a solution with $F > \dot{M}_{\rm b}$ (see Fig. 1). These may be genuinely unstable, leading to a common envelope, or actually lie in the region of thermal-timescale mass transfer marked in Fig. 2.

4. Stable mass transfer for large mass ratios

We have shown that mass transfer can be stable for large mass ratios M_2/M_1 provided that the conditions $D < 0, B < 0, M_1 \gtrsim 0.7M_\odot$ hold, and the system has a self-excited wind obeying Eq. (34). Observationally, these systems are characterized by episodic non-ejecting nova outbursts during which the white dwarf expands on average to a size comparable to the orbital separation: their average mass transfer rates are stably fixed very close to the threshold value $\dot{M}_{\rm b}$ (irradiation of the companion more efficient than for pure gravitational energy release, but less efficient than for the supersoft X–ray binaries). During hydrogen shell flashes these systems appear as very luminous (extreme–)ultraviolet sources, perhaps followed by a phase of quiet burning during which they may appear as a supersoft X-ray binary (cf. Shara et al. 1977; Prialnik et al. 1982). Between burning episodes the system will appear as a normal cataclysmic variable, but with an unusually large mass ratio. However the likely space density of these systems makes it improbable that any will be detected in this state.

The conditions Eq. (34) divide the M_1, M_2 plane as shown in Fig. 3. The evolution is therefore accessible to systems emerging from common-envelope evolution between these two curves. (Note that we do not consider the evolution of the binary before this point, so some of the mass ratios M_2/M_1 may not be realizable in practice, particularly where this quantity is large.) During the evolution the white dwarf mass increases, because the accreted matter is burnt without significant mass ejection (provided that the accreted matter is not lost during helium shell flashes, cf. Cassisi et al. 1998), but the total binary mass $M = M_1 + M_2$ decreases, as mass is lost in the self-excited wind. This kind of evolution thus ends either when mass transfer becomes formally unstable as the system crosses the right-hand boundary, or before this point if the white dwarf mass reaches the Chandrasekhar limit $\sim 1.4M_\odot$ first.

Since in general the mass loss timescale $M_2/(-\dot{M}_2) \gtrsim 10^8$ yr is longer than the thermal timescale of the companion we have assumed the thermal-equilibrium mass-radius exponent $\zeta \simeq 1$ in calculating the tracks shown in Fig. 3 for unevolved donor stars, and $\zeta = 0, -0.3$ (e.g. King et al. 1997a) in Fig. 4 for donors which are somewhat nuclear-evolved when mass transfer starts. However in a full calculation ζ should of course be adjusted self-consistently. As can be seen from the figures, the evolution is sensitive to the efficiency of wind driving through the factor $\chi = \eta_s^{1/2}\phi$. For large values (~ 1) of this parameter the wind losses are strong, so for unevolved donors (Fig. 3) an extreme mass ratio $M_2/M_1 \gg 1$ is needed to reduce B/D and thus the consequent mass transfer rate to $\sim \dot{M}_{\rm b}$ (cf. Eq. (33)). For small χ the required mass ratios are considerably lower. With evolved donors however (Fig. 4), high wind efficiencies are required for this evolution to be possible for any reasonable white dwarf masses at all.

Figs. 3 and 4 show that the white dwarf mass can grow quite substantially in this evolution (typically by $\sim 0.5M_\odot$). For unevolved donors and low wind efficiencies χ, or evolved donors and high wind efficiencies, it may even reach the Chandrasekhar limit, albeit from initial values which are themselves already quite large ($M_1 \gtrsim 1.2M_\odot$ in Fig. 3c) and/or require very massive secondary stars. However systems where M_1 has not reached $1.4M_\odot$ at the right-hand boundary may yet produce Type Ia supernovae, since on crossing this boundary they can become 'standard' supersoft X-ray binaries undergoing thermal-timescale mass transfer. For example, the track highlighted in Fig. 3b begins as a system with a white dwarf mass $M_1 = 0.78M_\odot$ and a companion mass $M_2 = 6M_\odot$, and reaches the rh boundary with $M_1 = 1.2M_\odot, M_2 = 3.75M_\odot$. From this point it is quite conceivable that the system can undergo thermal-timescale mass transfer and thus accrete the remaining $0.24M_\odot$ needed to produce a SN Ia.

5. Conclusions

We have considered the stability of mass transfer in binaries where self-excited wind losses are important. Mass transfer is stable for wind-driving laws slower than linear provided that the mass ratio M_2/M_1 is small. For wind-driving laws faster than linear, stability requires a large mass ratio. The linear case, considered by Ruderman et al. (1989) for low-mass X-ray binaries, is in practice indeterminate and fixed precisely by the deviations from this law: systems will settle in regimes where the wind law is faster or slower than linear depending on their mass ratios.

We have shown that in white dwarf binaries with large mass ratios M_2/M_1, self-excited winds naturally stabilize mass transfer at the threshold value $\dot{M}_{\rm b} \sim 10^{-8}M_\odot$ yr^{-1} for non-explosive nuclear burning of the accreted matter, as the wind-driving law is definitely faster than linear there. These systems undergo episodic non-ejecting nova outbursts with moderate expansion, and appear between the flashes as cataclysmic variables with an unusually large mass ratio. The white dwarf can accrete a substantial mass ($\sim 0.5M_\odot$), during this evolution, and may produce a Type Ia supernova directly in some cases. More probably, the descendants of this evolution are supersoft

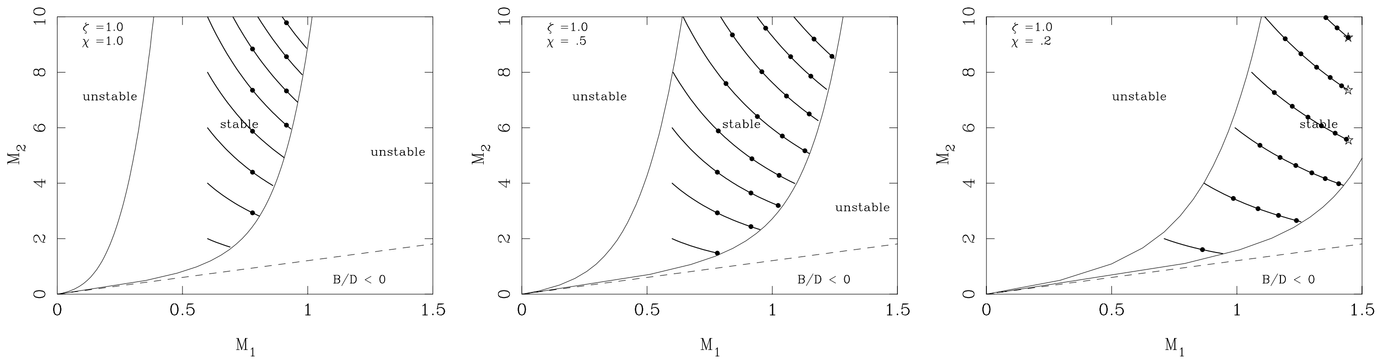

Fig. 3. Evolutionary tracks for high-mass wind-driven binaries in the M_1, M_2 plane for various values of the efficiency factor χ. The dots show intervals of 10^7 yr, and the stars denote systems where the white dwarf mass has reached the Chandrasekhar limit. Systems evolving across the right-hand boundary may continue mass transfer on a thermal timescale, allowing further efficient nuclear burning of the accreted matter. We assume unevolved donors losing mass on timescales longer than their thermal timescales, so that $\zeta \simeq 1$

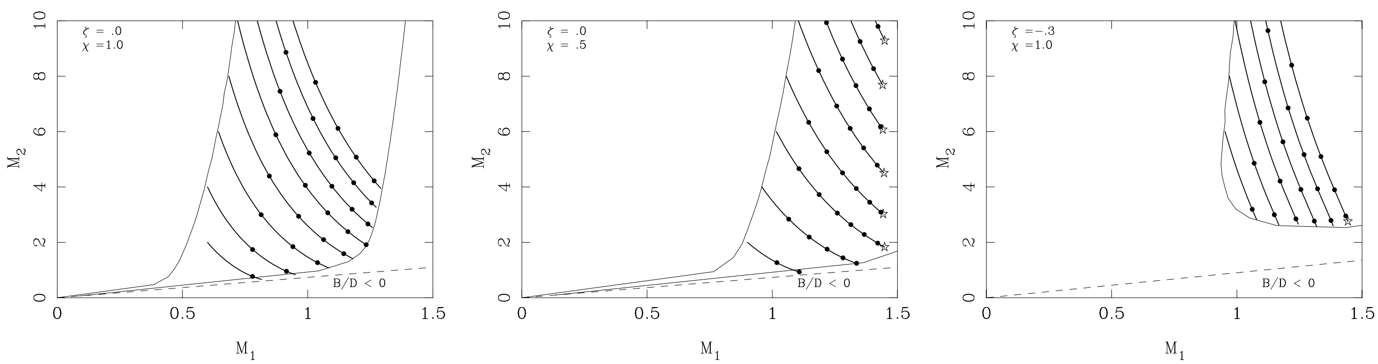

Fig. 4. Same as Fig. 3, but for evolved donors losing mass slowly, so that $\zeta \simeq 0, -0.3$

X-ray binaries with quite massive white dwarfs and moderate-mass companions, which then undergo thermal-timescale mass transfer and produce SNe Ia themselves.

It is known (e.g. Wheeler 1996; Livio 1996) that the usual thermal-timescale mass transfer picture has difficulty in producing the required rate of SNe Ia. The work of this paper has shown that the parameter space for SN Ia progenitors is probably considerably wider than hitherto thought, and that white dwarfs can accrete substantial amounts of mass non-explosively even with rather large mass ratios M_2/M_1. Full population synthesis calculations are needed to decide if these factors can resolve the current discrepancy between progenitor and supernova numbers.

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