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Wind-driven evolution of supersoft X-ray binaries with low-mass secondaries

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Abstract. We show that all supersoft X-ray binaries should excite strong winds ($\dot{M}_{\rm wind} \sim 10^{-7} M_\odot\,{\rm yr}^{-1}$) from the irradiated companion star. Scattering and absorption in this wind is likely to have a significant effect on the orbital X-ray light curve. In systems with a low-mass companion ($M_2/M_1 \lesssim 0.7$), such as the 4 hr supersoft X-ray binary 1E 0035.4-7230, the angular momentum loss in the wind may dominate the binary evolution and drive Roche lobe overflow at a rate comparable to the wind-loss rate. This may self-consistently sustain stable recurrent or steady-state hydrogen burning on the accreting white dwarf and keep the binary as a supersoft X-ray binary. These low-mass wind-driven supersoft X-ray binaries evolve towards longer periods, and can cross the period region populated by supersoft X-ray binaries driven by thermal-timescale mass transfer from a companion star more massive than the white dwarf. Starting from a massive white dwarf, low-mass wind-driven supersoft X-ray binaries may drive the white dwarf over the Chandrasekhar limit and produce a type Ia supernova.

Key words: accretion, accretion disks – stars: individual: 1E 0035.4-7230 – stars: mass-loss – novae, cataclysmic variables – supernovae: general – X-rays: stars

1. Introduction

It has been suggested that non-ejecting nova outbursts in systems with thermal-timescale mass transfer $\gtrsim 10^{-8} M_\odot\,{\rm yr}^{-1}$, may have long post-nova phases of stable burning of residual accreted hydrogen (e.g. Shara et al. 1977; Prialnik et al. 1982). Iben (1982) and Fujimoto (1982) argued that for even higher accretion rates $\gtrsim 10^{-7} M_1^{3.57} M_\odot\,{\rm yr}^{-1}$, the accreted hydrogen would burn in a steady state, i.e. at the accretion rate. At very high accretion rates $\gtrsim 8 \times 10^{-7} (M_1 - 0.5) M_\odot\,{\rm yr}^{-1}$, the white dwarf envelope expands on the horizontal shell-burning track.

After the discovery of several supersoft X-ray sources in the Magellanic Clouds by the EINSTEIN and ROSAT satellites (see review by Hasinger 1994), Van den Heuvel et al. (1992) used the model of steady-state nuclear burning on an accreting white dwarf with mass $M_1 \gtrsim 0.7 M_\odot$ to explain the observed characteristics of the binary supersoft X-ray sources CAL 83 and CAL 87. Optical observations have confirmed that about a dozen of the known supersoft X-ray sources are accreting binaries and there is an increasing amount of observational evidence that the compact accreting star in these supersoft X-ray binaries is a white dwarf with stable nuclear burning, although not necessarily in a steady state (see Van Teeseling 1998 for an observational overview). Van den Heuvel et al. argued that the required accretion rate for steady nuclear burning, without significantly expanding the white dwarf, could be obtained with thermal-timescale mass transfer from a near-main-sequence donor star with a mass in the range $\sim 1.3 - 2.5 M_\odot$. This is consistent with the observed orbital periods $\lesssim 1$ day for most of the known supersoft X-ray binaries. The 4 hr supersoft X-ray binary 1E 0035.4-7230 is clearly a supersoft X-ray binary of a different kind.

Van den Heuvel et al. also noted that their model implied the presence of strong winds from the heated side of the companion and the disk. Evolutionary calculations of supersoft X-ray binaries have been performed by Rappaport et al. (1994), Yungelson et al. (1996), and DiStefano & Nelson (1996). Only the latter included the evolutionary effects of mass loss from the white dwarf and the secondary, but the amount of mass loss was an unknown free parameter. In this paper, we present an estimate of the amount of mass loss from the irradiated secondary and discuss the evolutionary consequences of very strong mass loss from supersoft X-ray binaries with a companion star less massive than the white dwarf. As an example of such a system, we discuss the 4 hr supersoft X-ray binary 1E 0035.4-7230 in more detail.

2. The 4.1 hr supersoft X-ray binary 1E 0035.4-7230

Among the supersoft X-ray binaries with known orbital periods, 1E 0035.4-7230 (= SMC 13) has, with 4.126 hr, by far the lowest period (Schmidtke et al. 1996). This period is significantly lower than that required for a binary with thermal-timescale mass transfer with a rate of $\gtrsim 4 \times 10^{-8} M_\odot\,{\rm yr}^{-1}$, necessary for steady-state burning (cf. Rappaport et al. 1994). With a white dwarf mass $> 0.6 M_\odot$, and a secondary on the main-sequence (or under-massive for its radius), a 4 hr orbital period implies a mass ratio $M_2/M_1 \lesssim 0.7$. Clearly, 1E 0035.4-7230 does not have steady-state burning or there is a mechanism which increases the mass transfer rate in this system to a steady-state burning value.

Kahabka & Ergma (1997) have proposed that 1E 0035.4-7230 is a cataclysmic variable, currently in a phase of residual hydrogen burning after a mild shell flash. Necessary conditions to have such a long phase of residual hydrogen burning (at least 28 years since its discovery with EINSTEIN; Seward & Mitchell 1981), would be low CO abundances of the burning matter and a hot low-mass ($\sim 0.6 - 0.7 M_{\odot}$) white dwarf. The low CO abundances require low metal abundances of the accreted matter and a thick helium layer on top of the CO core to prevent the mixing of accreted matter with the core. Due to compressional heating between previous nova outbursts, the white dwarf would have reached the hot equilibrium temperature already in a few 10^6 yrs after the onset of mass transfer. Also Sion & Starrfield (1994) note the possible existence of extremely hot, low-mass white dwarfs powered by stable hydrogen burning for thousands of years while accreting at $\sim 10^{-8} M_{\odot}$ yr^{-1}.

If the phase of residual hydrogen burning in 1E 0035.4-7230 lasts only several hundred years, as stated by Kahabka & Ergma (1997), and the recurrence time of the nova outbursts were $\sim 10^5$ yrs, then only a fraction of 0.01–0.001 of these binaries would appear at any time as a supersoft X-ray source. As the accretion rate has to be $> 10^{-9} M_{\odot}$ yr^{-1}, only novalike variables above the period gap could turn into such recurrent supersoft X-ray binaries. Scaling the estimated number of novalike variables in the Milky Way (e.g. Kolb 1996) with the masses of the SMC and the Milky Way gives a total of $\sim 10^3$ novalike variables in the SMC. If we take into account that perhaps only a fraction of $\sim 0.1 - 0.4$ of the supersoft X-ray binaries in the SMC would be detectable with ROSAT because of interstellar absorption (Rappaport et al. 1994), this would imply that practically all novalike variables in the SMC would have to be like 1E 0035.4-7230. Because normal novae are also observed in the SMC (see Welch 1997 for the most recent one), this cannot be true. We conclude that it is likely that the supersoft X-ray phase of 1E 0035.4-7230 lasts significantly longer than several hundred years.

The accreting object in 1E 0035.4-7230 has a luminosity of the order of 10^{37} erg s^{-1} (with an uncertainty of at least a factor of 3) and a temperature of $\sim 5 \times 10^5$ K consistent with a hot white dwarf (Van Teeseling et al. 1996a). The optical flux of 1E 0035.4-7230 is dominated by the strongly irradiated accretion disk and secondary. The changing aspect of the irradiated secondary causes the observed $\Delta V \sim 0.3$ mag optical modulation. Because the optical flux is dominated by reprocessed soft X-rays, there is no observational clue about the accretion rate, and there is, therefore, no observational evidence that the system accretes below (or in) the steady-state shell burning regime. Independently of how 1E 0035.4-7230 became a supersoft X-ray binary, in the next sections we show that the strong (and observed) irradiation of the companion star will inevitably lead to a very strong wind from its heated side. If the supersoft X-ray phase lasts longer than $\sim 10^3$ yrs, this wind may dominate the binary evolution, increase the mass accretion rate and force the binary to stay in the steady-state nuclear burning regime.

3. Winds in supersoft X-ray binaries

3.1. Observational evidence

The Balmer emission lines in the spectra of RX J0019.8+2156 and RX J0513.9-6951 show P Cygni absorption which indicates the presence of a wind reaching velocities exceeding 1000 km s^{-1} (Beuermann et al. 1995; Crampton et al. 1996; Reinsch et al. 1996). The ultraviolet spectra of RX J0513.9-6951 and CAL 83 show evidence for similar high-velocity outflows (Gänsicke et al. 1998). Unfortunately, the published optical spectra of 1E 0035.4-7230 have too low signal-to-noise to detect similar P Cygni absorption profiles (Van Teeseling et al. 1996b; Crampton et al. 1997). Part of the strong He II emission line, which has been observed in all supersoft X-ray binaries, may also originate in a wind, similar as in Wolf-Rayet stars. Indeed, its equivalent width increases during primary eclipse in CAL 87 suggesting a large extent of the line-emitting region (Van den Heuvel et al. 1992).

Another observational feature which may be at least partly due to scattering in a wind from the binary is the presence of X-ray orbital modulation in 1E 0035.4-7230, RX J0019.8+2156, and CAL 87 (Kahabka 1996; Crampton et al. 1997). Because the scattering must be asymmetric with respect to the X-ray source, a large fraction of this wind probably originates on the companion star, as is confirmed by the fact that in all systems X-ray minimum occurs when the companion star is closest to us.

We can obtain an order of magnitude estimate of the wind mass loss rate in 1E 0035.4-7230 by assuming a constant-velocity spherically symmetric wind from the companion star with a density profile $\rho = \rho_{\rm R}(R^2/r^2)$, with R the radius of the companion star. Then the observed column density is $\sim \rho_{\rm R} R$, with most of the scattering matter within the binary system. An observed X-ray modulation of a factor of ~ 2 requires a scattering optical depth close to unity, equivalent to a column density of ~ 3 g cm^{-2}. Using a constant wind velocity of the order of the sound velocity $v_{\rm s} \sim 10^7$ cm s^{-1}, the resulting wind mass loss rate from the companion is

$$\dot{M}_{\rm w2} \sim -2\pi R^2 \rho_{\rm R} v \sim -6\pi R v_{\rm s} \sim -10^{-7}\ M_{\odot}{\rm yr}^{-1}. \qquad (1)$$

3.2. An irradiation-induced wind from the companion star

Strong irradiation of the companion star by a luminous X-ray source will form an extended corona on top of the companion star atmosphere and induce a strong stellar wind (Basko & Sunyaev 1973). The very soft nature of the irradiating X-rays in supersoft X-ray binaries is even more effective in driving a stellar wind, because most of the energy is absorbed above the Rosseland photosphere (mainly by He II), in contrast to hard X-ray irradiation, where most of the energy will be absorbed below the photosphere. Since supersoft X-ray binaries have a very high efficiency of energy release per accreted gram due to nuclear shell burning, giving a luminosity close to the Eddington limit, it is clear that there are no binaries in nature in which the production of an irradiation-induced stellar wind can be as efficient as in these sources.

The mass outflow rate is essentially determined by the product of the density and the outflow velocity at the base of the corona (cf. Eq. (1)). Basko & Sunyaev (1973) show that there are two regimes in driving a stellar wind by irradiation: for low X-ray fluxes the mass outflow rate is directly proportional to it. Increasing the X-ray irradiation shifts the temperature jump deeper into the atmosphere with a higher density and the outflow rate increases. For high X-ray fluxes, recombination and optical depth effects in the corona will be important and the density at the jump will be proportional only to the square root of the flux. It is easy to verify that in the case of supersoft X-ray binaries, we are well within this high-flux regime. A simple way to describe this regime is to consider the corona as a Strömgren zone. Then the number flux of ionizing photons n_i will be equal to the recombination rate. We approximate the Strömgen corona as an isothermal layer with an exponential density structure and a scale height $H = v_s^2 R_2^2/GM_2$, and assume that He II is the main absorber ($h\nu_{\rm T} = 54.4\,$eV). Then

$$n_i = \int_{\nu_{\rm T}}^{\infty} \frac{F_\nu}{h\nu} {\rm d}\nu = \alpha({\rm He}^+, T) \int N_{\rm e} N_{{\rm He}^+} {\rm d}z \qquad (2)$$

which gives with $N_{{\rm He}^+} = x_{\rm He} N_{\rm e} = x_{\rm He} N_0 \exp(-z/H)$,

$$n_i = \alpha N_0^2 x_{\rm He} \frac{H}{2}. \qquad (3)$$

Note that for high irradiation fluxes, Eq. (31) of Basko & Sunyaev (1973) for the density at the base of the corona reduces to our Eq. (3). The wind mass loss rate from the irradiated hemisphere of the companion star is

$$\dot{M}_{\rm w2} \simeq -\phi \pi R_2^2 N_0 \mu m_{\rm H} u v_{\rm s}, \qquad (4)$$

where u is the isothermal Mach number immediately above the temperature inversion, i.e. the ratio of the velocity to the isothermal sound velocity $v_{\rm s}$. 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. Shadowing by the accretion disk may reduce ϕ, but observations show that at least in the case of 1E 0035.4-7230 this effect is probably not very large (Van Teeseling et al. 1998). Eliminating N_0 in this equation using Eq. (3) gives

$$\dot{M}_{\rm w2} \simeq -\phi \pi R_2 \sqrt{\frac{2 n_i G M_2}{\alpha x_{\rm He}}} \mu m_{\rm H} u. \qquad (5)$$

According to Basko et al. (1977), the flow immediately above the temperature jump is still subsonic. However, it cannot be very subsonic, because then the flow would resemble a quasi-static corona, which would be unstable: numerical calculations indicate that already for soft X-ray irradiation with very low fluxes no hydrostatic equilibrium can be achieved (Hessman et al. 1997). Another reason why the flow cannot be very subsonic above the temperature inversion is the Roche-lobe potential: the flow must become supersonic before it reaches the maximum in the gravitational potential. Therefore, the flow must become supersonic at the inner Lagrange point and within about two stellar radii elsewhere in the equatorial plane (corresponding to $u \gtrsim 0.3$). At higher latitudes the velocity at the base of the corona could be lower, leading to an outflow concentrated towards the equatorial plane. Unfortunately, very little is known about the non-hydrostatic structure of an atmosphere with extremely strong irradiation by a luminous supersoft X-ray source and the acceleration in the temperature jump, and we will adopt $u = 0.3$ in the remaining discussion.

The temperature of the corona, and therefore $\alpha({\rm He}^+, T)$, will also depend on the X-ray flux. The temperature will range from approximately the ionization temperature of He II ($\sim 3 \times 10^4$ K) for low irradiation fluxes to almost the temperature of the irradiation for very high irradiation fluxes and a subsonic outflow. We will use $T \sim 5 \times 10^5$ K and $\alpha({\rm He}^+, T) \sim 1 \times 10^{-13}$ cm^3 s^{-1}. We take $x_{\rm He} = 0.08$, $\mu m_{\rm H} = 3 \times 10^{-24}$ g, and

$$n_i \simeq \frac{3 \times 10^{23}}{a_{11}^2} L_{37} \eta_s. \qquad (6)$$

Here η_s measures the efficiency of the primary's spectrum in producing ionizing photons and driving a wind, where for supersoft X-ray irradiation temperatures of a few 10^5 K, $\eta_s \simeq 1$. L_{37} is the luminosity of the X-ray source in 10^{37} erg s^{-1}. Then we find

$$\dot{M}_{\rm w2} \simeq -3 \times 10^{-7} \phi \frac{r_2}{a_{11}} (m_2 L_{37} \eta_s)^{1/2}\ M_\odot {\rm yr}^{-1}, \qquad (7)$$

where a_{11} is the orbital separation in 10^{11} cm, and r_2 and m_2 the radius and mass of the companion star in solar units. It should be noted that for lower irradiation temperatures $T \lesssim 10^5$ K, and with hydrogen as the dominant absorber, the constant in Eq. (7) would be almost identical, implying an effective $\eta_s \simeq 1$.

For an accretion rate $\dot{m} \gtrsim 1.32 m_1^{3.57}$, with $\dot{m}$ in units of $10^{-7} M_\odot$ yr^{-1} and m_1 the white dwarf mass in $M_\odot$, the luminosity of the white dwarf comes from steady-state shell burning of the accreted hydrogen, and we can replace the luminosity by the accretion rate using $L_{37} \simeq 2.9\dot{m}$ (Iben 1982). The same relation holds for the long-term average luminosity of supersoft X-ray binaries with recurrent stable hydrogen burning, provided that during the shell flash no common envelope is formed and no significant mass is lost from the white dwarf envelope, which is roughly true for an accretion rate (see Fig. 9 in Fujimoto 1982)

$$\dot{m} \gtrsim 0.63 \times 10^{-0.8 m_1}. \qquad (8)$$

More generally we can write (neglecting a wind from the white dwarf)

$$L_{37} \simeq 2.9 \dot{m} \eta_a, \qquad (9)$$

where η_a measures the luminosity per gram of matter accreted relative to the value for nuclear shell burning. Eq. (7) becomes

$$\dot{M}_{\rm w2} \simeq -5 \times 10^{-7} \phi \frac{r_2}{a_{11}} (m_2 \dot{m} \eta_s \eta_a)^{1/2}\ M_\odot {\rm yr}^{-1}. \qquad (10)$$

In practice, the wind mass loss rate may not reach the value in Eq. (10). Saturation by an increasing amount of photoabsorption has already been included by approximating the outflowing matter as a Strömgren corona, and causes the wind mass loss rate to be proportional only to the square root of the irradiating flux (Eq. (7)). Our simple one-dimensional treatment breaks down, however, when scattering between the X-ray source and the companion star significantly reduces the amount of irradiation which is absorbed in the outflowing wind. In the worst case (when the wind is almost fully ionized) this will happen when (cf. Eq. (1))

$$\dot{M}_{\rm w2,max} \simeq -3\tau_{\rm e}\pi R_2\sqrt{\frac{2GM_2}{R_2}} \sim -6\times 10^{-7}\sqrt{m_2 r_2}\, M_{\odot}{\rm yr}^{-1}. \quad (11)$$

However, if also a wind from the white dwarf and/or the accretion disk is present, Eq. (11) is an overestimate. Comparing Eqs. (7) and (11), we see that the wind saturates when

$$L_{37} \gtrsim \frac{a_{11}^2}{r_2\eta_s\phi^2}. \quad (12)$$

It appears that most supersoft X-ray binaries, with a luminosity of a few times 10^{37} erg s^{-1}, are very close to this saturation limit (note that 1E 0035.4-7230 gives $|\dot{M}_{\rm w2}| \sim 1\times 10^{-7} M_{\odot}$ yr^{-1}, which is consistent with the estimate from the X-ray modulation in Eq. (1)). Therefore, it seems unavoidable that in supersoft X-ray binaries the companion star loses mass with a rate of $10^{-7} - 10^{-6} M_{\odot}$ yr^{-1}. This must have observational effects on the orbital X-ray light curves in the smaller systems (as observed in 1E 0035.4-7230, RX J0019.8+2156, and CAL 87). However, we note that a final statement about the exact amount of saturation and the effect of the wind on the orbital light curves requires a full 3-D hydrodynamical radiative-transport calculation. In the next section, we will show that such a strong wind may also have a significant effect on the binary evolution of supersoft X-ray binaries with a low-mass companion, like 1E 0035.4-7230.

4. Winds and binary evolution

The mass and particularly angular momentum carried off by a strong wind from either component of a binary system must affect the orbital elements and thus the long-term evolution of the binary. Tout & Hall (1991) considered the effects of the stellar wind of the mass-losing star in Algol systems, and concluded that this could drive the binary evolution in some cases. King & Kolb (1995) discussed the case in which the mass-gaining component has a wind loss rate proportional to the mass transfer rate, and recently Li & Van den Heuvel (1997) have studied the effect of wind losses from the white dwarf in stabilizing what would otherwise be thermal-timescale mass transfer in supersoft X-ray binaries. Here we are mainly concerned with a case not considered by Tout & Hall (1991), namely where the wind from the mass-losing star is driven by irradiation by the primary, and thus itself a consequence of the mass transfer. This situation can obviously occur quite independently of the evolutionary state of the donor star. The short orbital period of 1E0035.4-7230 suggests that the donor is not nuclear-evolved, but there is no reason to exclude the possibility of evolved donors from our general discussion of wind-driven evolution.

To see how the various cases fit together we first consider the general situation in which the two stars have wind mass-loss rates $\dot{M}_{\rm w1}, \dot{M}_{\rm w2} < 0$ respectively, angular momentum can be lost by the binary in other ways, and the secondary star may expand or contract for reasons other than mass loss. We parameterize the associated angular momentum loss as β_1, β_2 times the specific orbital angular momenta $j_1 = M_2J/M_1M, j_2 = M_1J/M_2M$, where M and J are the total binary mass and orbital angular momentum, i.e.

$$J = M_1M_2\left(\frac{Ga}{M}\right)^{1/2}. \quad (13)$$

The quantity controlling the evolution is the radius $R_{\rm L}$ of the Roche lobe of the mass-losing (secondary) star, for low mass ratios $M_2/M_1 \lesssim 0.8$ given approximately by

$$\frac{R_{\rm L}}{a} = 0.462\left(\frac{M_2}{M}\right)^{1/3} \quad (14)$$

(Paczyński 1971). Logarithmically differentiating and combining these two equations gives

$$\frac{\dot{R}_{\rm L}}{R_{\rm L}} = \frac{2\dot{J}}{J} - \frac{2\dot{M}_1}{M_1} - \frac{5}{3}\frac{\dot{M}_2}{M_2} + \frac{2}{3}\frac{\dot{M}}{M}. \quad (15)$$

The derivatives appearing on the rhs of this equation are given by

$$\dot{M}_1 = -\dot{M}_{\rm tr} + \dot{M}_{\rm w1} \quad (16)$$

$$\dot{M}_2 = \dot{M}_{\rm tr} + \dot{M}_{\rm w2} \quad (17)$$

$$\dot{M} = \dot{M}_{\rm w1} + \dot{M}_{\rm w2} \quad (18)$$

$$\dot{J} = \beta_1\dot{M}_{\rm w1}j_1 + \beta_2\dot{M}_{\rm w2}j_2 + \dot{J}_{\rm sys}, \quad (19)$$

where $\dot{M}_{\rm tr} < 0$ is the mass transfer rate from the secondary to the primary, the terms in j_1, j_2 are the wind angular momentum loss rates, and $\dot{J}_{\rm sys}$ denote all other forms of angular momentum loss from the binary orbit, such as magnetic braking or gravitational radiation. Combining with Eq. (15) and the definitions of j_1, j_2 gives

$$\frac{\dot{R}_{\rm L}}{R_{\rm L}} = -\frac{2\dot{M}_{\rm tr}}{M_2}\left[\frac{5}{6} - \frac{M_2}{M_1}\right] + \frac{\dot{M}_{\rm w1}}{MM_1}\left[2(\beta_1-1)M_2 - \frac{4}{3}M_1\right] + \frac{\dot{M}_{\rm w2}}{MM_2}\left[\left(2\beta_2 - \frac{5}{3}\right)M_1 - M_2\right] + \frac{2\dot{J}_{\rm sys}}{J}. \quad (20)$$

To see the effect of this change of $R_{\rm L}$ we must compare it with the change of the secondary's radius R_2 on mass loss; we write $\dot{R}_2/R_2 = \dot{R}_s/R_s$ if the star changes its radius independently of mass loss, e.g. via thermal expansion across the Hertzsprung gap or nuclear expansion, and

$$\frac{\dot{R}_2}{R_2} = \zeta\frac{\dot{M}_2}{M_2} = \zeta\frac{\dot{M}_{\rm tr}}{M_2} + \zeta\frac{\dot{M}_{\rm w2}}{M_2}, \quad (21)$$

if the radius change is proportional to the mass loss rate. Here ζ is the *effective* mass-radius index, i.e. the mass–radius index actually followed by the secondary along its evolutionary track. The index ζ therefore depends both on the nature of the secondary and the rate at which it is losing mass, and must in general be calculated self-consistently through the evolution. However we can infer its likely value quite easily in many cases. For example, a low-mass ($M_2 \lesssim 0.6 M_\odot$) unevolved star expands ($\zeta \simeq -1/3$) if it loses mass adiabatically, i.e. on a timescale short compared with its thermal timescale, but contracts and remains close to the main sequence ($\zeta \simeq 1$) if the mass-loss timescale is longer than its thermal time. A somewhat more massive ($0.6 M_\odot \lesssim M_2 \lesssim 2 M_\odot$) unevolved star shrinks rapidly ($\zeta \gg 1$) on adiabatic mass loss, but more gently ($\zeta \simeq 1$) if the mass-loss timescale is longer than thermal. For a system with small mass ratio like 1E 0035.4-7230, rapid adiabatic mass-loss from the low-mass secondary with a deep convective envelope implies $\zeta = -1/3$ to a good approximation.

For a given binary, it may happen that the difference $R_{\rm L} - R_2$ decreases with time. Then the mass transfer rate will grow exponentially on a timescale $H/R_2 \sim 10^{-4}$ times the current mass transfer timescale, where H is the scale height of the secondary near the inner Lagrange point (Ritter 1988). If on the other hand, a stable state exists in which the Roche lobe and stellar radius can move in step, the binary will rapidly find this state. The condition $\dot R_{\rm L}/R_{\rm L} = \dot R_2/R_2$ then determines the mass transfer rate $\dot M_{\rm tr}$. Combining Eqs. (20) and 21 gives

$$\frac{\dot M_{\rm tr}}{M_2}\left[\frac{\zeta}{2} + \frac{5}{6} - \frac{M_2}{M_1}\right] = \frac{\dot M_{\rm w1}}{M M_1}\left[(\beta_1 - 1)M_2 - \frac{2}{3}M_1\right] + \frac{\dot M_{\rm w2}}{M M_2}\left[\left(\beta_2 - \frac{5}{6} - \frac{\zeta}{2}\right)M_1 - \left(\frac{1}{2} + \frac{\zeta}{2}\right)M_2\right] + \frac{\dot J_{\rm sys}}{J} - \frac{\dot R_{\rm s}}{2R}, \tag{22}$$

where one of ζ or $\dot R_s$ will be zero in any given case. Eq. (22) agrees in the appropriate limits with the equations of King & Kolb (1995) and Tout & Hall (1991), and with earlier standard treatments. For example, for cataclysmic variables the wind and stellar expansion terms are negligible, and we get

$$\frac{\dot M_{\rm tr}}{M_2} = \frac{\dot J_{\rm sys}/J}{\zeta/2 + 5/6 - M_2/M_1}, \tag{23}$$

giving the familiar results

$$\frac{\dot M_{\rm tr}}{M_2} = \frac{\dot J_{\rm sys}/J}{2/3 - M_2/M_1}, \tag{24}$$

$$\frac{\dot M_{\rm tr}}{M_2} = \frac{\dot J_{\rm sys}/J}{4/3 - M_2/M_1}, \tag{25}$$

in the cases where $\dot J_{\rm sys}/J$ drives adiabatic ($\zeta = -1/3$) and slow ($\zeta = 1$) mass transfer respectively. For mass transfer driven by radius evolution one neglects all the angular momentum loss terms and sets $\zeta = 0$, obtaining the familiar

$$\frac{\dot M_{\rm tr}}{M_2} = \frac{-\dot R_{\rm s}/2R_2}{5/6 - M_2/M_1}. \tag{26}$$

With $\dot R_s/R_2 \sim 1/t_{\rm nuc}$, this equation describes mass transfer driven by the nuclear expansion of the secondary. Similarly a secondary crossing the Hertzsprung gap expands on a thermal timescale, so $\dot R_s/R_2 \sim 1/t_{\rm KH} > 0$. The thermal-timescale mass transfer often invoked (Van den Heuvel et al. 1992) for driving the supersoft X-ray binaries corresponds to $\dot R_s/R_2 \sim -1/t_{\rm KH} < 0$, i.e. the donor star would like to *shrink* on a thermal timescale in response to mass loss. Stars reasonably close to the main sequence have this property. Clearly mass transfer would stop fairly quickly unless the Roche lobe also shrank, so this type of evolution requires a mass ratio $M_2/M_1 \gtrsim 1$. Mass transfer is thus formally unstable on a thermal timescale. Dynamical–timescale stability is nevertheless assured since near main-sequence or slightly evolved secondaries with masses $\gtrsim 0.8 M_\odot$ have $\zeta = \zeta_{\rm ad} \gg 1$ for rapid (adiabatic) mass loss (it is consistent to define an adiabatic mass–radius index since the stars are assumed to be in hydrostatic equilibrium). Thus the star and Roche lobe shrink stably together on a thermal timescale. To our knowledge, there is so far no published calculation following this type of evolution in detail.

We note finally that we recover Tout & Hall's requirement $\zeta < 1/3$ for mass transfer to be driven entirely by wind losses from the secondary (all terms on rhs of Eq. (22) neglected except for that in $\dot M_{\rm w2}$, with $\beta_2 = 1$). This shows that wind-driven mass loss is impossible if the secondary can shrink rapidly enough on mass loss (cf. Fig. 1).

In this paper, we are mainly interested in the case of strong wind loss from the secondary and so assume $\beta_2 = 1$ and $\dot M_{w1} = 0$ in the remaining discussion. We assume also that the wind mass loss timescale $M_2/\dot M_{w2}$ is short compared with that for thermal or nuclear radius evolution or systemic angular momentum loss, and so neglect the terms in $\dot R_{\rm s}$, $\dot J_{\rm sys}$ in Eq. (22). This means that the mass loss is essentially adiabatic, and we can use values $\zeta \simeq \zeta_{\rm ad} \sim -1/3$ for the low-mass secondaries we shall consider.

Fig. 1 illustrates that there are two regions in the $q = M_2/M_1$ vs. ζ plane for which a wind from the secondary will tend to drive a positive mass transfer rate. It is important to note that ζ is the *effective* mass-radius exponent, and therefore depends on the rate of mass loss as well as the structure of the star. Systems may therefore adjust the value of ζ in the course of their evolution, and indeed this is what the donors in thermal-timescale mass-transfer systems do.

For an irradiation-induced wind loss rate proportional to $\dot M_{\rm tr}^{1/2}$ (see Eq. (10)), only the solutions with

$$q < \frac{1 - 3\zeta}{3(1+\zeta)} \quad \& \quad q < \frac{\zeta}{2} + \frac{5}{6} \tag{27}$$

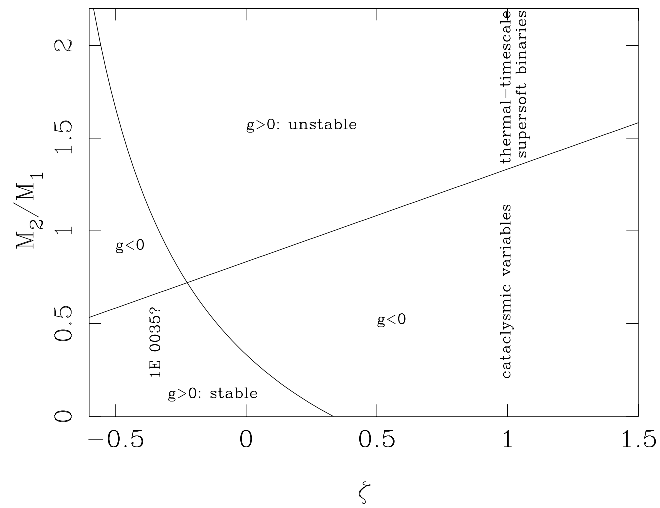

Fig. 1. The regions in the M_2/M_1 vs. effective ζ plane where a wind from the secondary drives mass transfer ($g > 0$; see Eq. (29)) or stabilizes mass transfer ($g < 0$). The terms *stable* and *unstable* refer to the stability of the solutions with $m_{w2} \propto \dot{m}^{1/2}$. Also indicated are the positions of cataclysmic variables, the possible position of 1E 0035.4-7230, and supersoft X-ray binaries with thermal timescale mass transfer. The latter are arbitrarily assigned effective mass-radius indices $\zeta \gtrsim 1$ since they shrink on mass loss

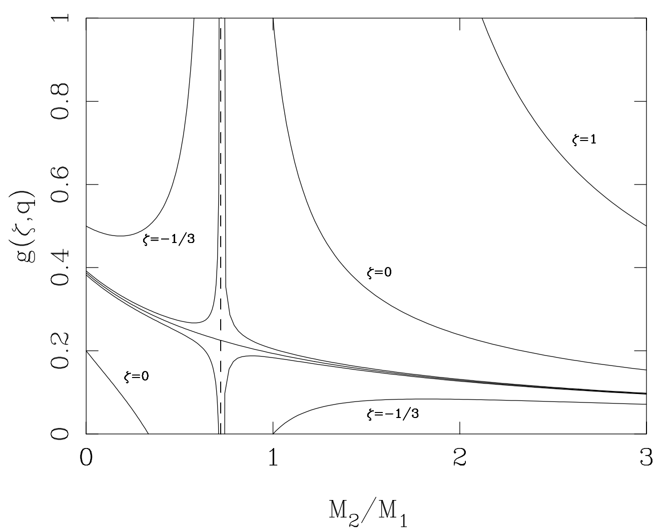

Fig. 2. The values of $g(\zeta, q)$ as a function of the mass ratio for different values of the effective mass-radius index ζ. Note that for mass ratios $\gtrsim 0.8$, the used approximation for the Roche lobe (Eq. (14)) is not very accurate, and that this part is only shown for illustrative purposes

are stable. For $\zeta = -1/3$, stable wind-driven mass transfer is possible for any mass ratio $q < 2/3$, as is likely to be true for 1E 0035.4-7230.

More generally, mass transfer driving ($\dot{M}_{\rm tr} < 0$) requires that the primary wind should be sufficiently weak compared with that from the secondary, e.g. for $q < \zeta/2 + 5/6$ and $\beta_1 = \beta_2 = 1$,

$$\frac{\dot{M}_{\rm w1}}{\dot{M}_{\rm w2}} < \frac{3}{4q}\left[\frac{1}{3} - \zeta - (1+\zeta)q\right]. \qquad (28)$$

For $\zeta = -1/3$, $\dot{M}_{\rm w1}/\dot{M}_{\rm w2} < 1/4$ would imply mass transfer driving for any $q < 2/3$. If the luminosity of the shell-burning white dwarf is well below the Eddington luminosity, one would indeed expect the irradiation-driven wind from the secondary to be stronger than that from the primary plus the irradiated accretion disk, since the radiation is incident normally on much of the secondary, and its surface gravity is probably lower than that of the inner regions of the accretion disk. Evidence for this in the case of the supersoft X-ray binaries is the fact that some of these systems show a large orbital modulation in their X-ray light curves, possibly due to scattering in an asymmetric wind. A dominant wind from the primary component would be axisymmetric and could not show such an orbital effect.

Accordingly in the following we consider the case $\dot{M}_{w1} = 0$, $\beta_2 = 1$, and drop the label '2' from $\dot{M}_{\rm w2}$. Thus, Eq. (22) for the mass transfer rate reduces to

$$\dot{M}_{\rm tr} = \dot{M}_{\rm w} g(\zeta, q), \qquad (29)$$

where

$$g(\zeta, q) = \frac{(1-3\zeta) - 3(1+\zeta)q}{(1+q)(5+3\zeta-6q)}. \qquad (30)$$

Fig. 2 shows the positive values of g as a function of q for different values of ζ. For small mass ratios and $\zeta = -1/3$ we find $g \simeq 0.5$, except for mass ratios very close to the stability limit $q = 2/3$. For such systems, we can thus draw the general conclusion that wind losses typically drive Roche lobe overflow of order one-half of the wind loss rate.

5. Low-mass supersoft X-ray binaries driven by self-excited winds

We can now combine the results of the last two sections, i.e. consider the case where the primary's luminosity excites a wind which drives Roche lobe overflow and thus check if this can sustain the primary's luminosity. Ruderman et al. (1989) attempted to demonstrate such a 'bootstrap' evolution for low-mass X-ray binaries. Here, we show that some supersoft X-ray binaries are more promising candidates for this.

We first combine Eq. (10) with the Roche lobe condition Eq. (14) to replace the ratio r_2/a_{11}, which gives

$$\dot{m}_{\rm w} = 3.5\phi m_2^{5/6} m^{-1/3} (\eta_s \eta_a \dot{m})^{1/2}, \qquad (31)$$

where $\dot{m}, \dot{m}_{\rm w}$ are $-\dot{M}_{\rm tr}, -\dot{M}_{\rm w}$ in units of $10^{-7} M_\odot$ yr^{-1}. If the mass transfer is driven by a wind from the secondary, so that $\dot{m}_{\rm w} = \dot{m}/g$ (Eq. (29)), we get

$$\dot{m} = 0.26\phi^2 \left(\frac{m_2}{0.1}\right)^{5/3} m^{-2/3} \eta_s \eta_a g^2. \qquad (32)$$

This is only self-consistent for a low-mass wind-driven supersoft X-ray binary (for which we have normalized $\eta_s \simeq \eta_a \simeq 1$)

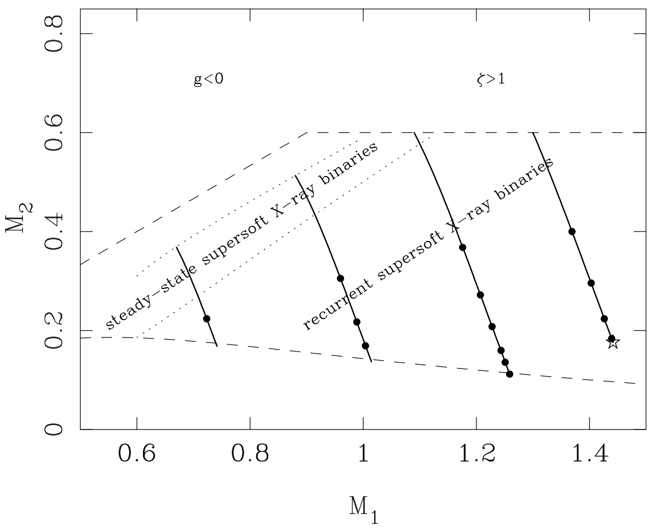

Fig. 3. The upper dashed line corresponds to the maximum mass of a low-mass secondary with $\zeta = -1/3$ for which a wind drives mass transfer (i.e. $q < 2/3$ and $M_2 < 0.6M_\odot$). The upper dotted line corresponds to the maximum accretion rate for which steady-state hydrogen shell burning is possible (core mass–luminosity relation for red giants). The lower dashed line corresponds to an accretion rate at which the white dwarf expands significantly (to $1R_\odot$) during recurrent hydrogen shell burning (Fujimoto 1982) and irradiation-induced winds become impossible. The lower dotted line corresponds to the stability line of Iben (1982) which gives the lowest possible accretion rate for stable steady-state hydrogen burning. Wind-driven evolution of supersoft X-ray binaries with a low-mass secondary is possible in the region bounded by the 'red-giant' line, the $M_2 = 0.6M_\odot$, and the lower dashed line. The solid lines are evolutionary tracks with the dots marking 10^6 yrs since the start on the 'red-giant' line or the $M_2 = 0.6M_\odot$ line

with $\dot{m}$ in the permitted range for stable nuclear shell burning without the formation of a common envelope (see Eq. (8)). Combining Eqs. (8) and (32), we get the requirement

$$m_2 \gtrsim 0.17 m^{2/5} (\eta_s \eta_a)^{-3/5} (g\phi)^{-6/5}\, 10^{-0.48 m_1}. \tag{33}$$

Hence, self-consistent wind-driven evolution is possible for low-mass supersoft X-ray binaries obeying (33). Fig. 3 shows the area of the M_1, M_2 plane where this evolution is permitted, for the case $\eta_s = \eta_a = \phi = 1, \zeta = -1/3$. In calculating the accretion and wind loss rates for Fig 3, we have taken into account that the wind becomes saturated for high accretion rates (in the 'steady-state strip') and that the wind loss rate cannot exceed $\dot{M}_{\rm w2,max}$ (Eq. (11)).

Combining the Roche lobe condition Eq. (14) with Kepler's 3rd law shows that the binary period P obeys

$$\frac{\dot{P}}{P} = \frac{3\zeta - 1}{2} \frac{\dot{M}_2}{M_2}. \tag{34}$$

Since as stated above we generally have $\zeta = -1/3$ for this type of evolution, Eq. (34) implies that the binary period evolves as

$$P \propto M_2^{-1}. \tag{35}$$

The constant of proportionality here is set by the initial condition, i.e. the orbital period and secondary mass at which the wind-driven evolution started. If at this time the secondary was close to the main sequence with initial mass $m_{\rm i} M_\odot$, its initial period would have been

$$P_{\rm i} \simeq 9 m_{\rm i}\ {\rm hr}. \tag{36}$$

Hence, the period evolution of such a system is described by

$$P \simeq 0.11 \frac{P_{\rm i}^2}{m_2}\ {\rm hr}. \tag{37}$$

If 1E 0035.4-7230 has just started wind-driven evolution, its final period could be as long as ~ 10 hr with a $\sim 0.8M_\odot$ white dwarf and up to ~ 18 hr with a very massive white dwarf. The total possible period range of low-mass wind-driven supersoft X-ray binaries (with an unevolved companion) is $\sim 2-30$ hr. The wind-driven systems may therefore overlap the period range of the high-mass supersoft X-ray binaries with thermal-timescale mass transfer, or even evolve beyond it. The evolutionary lifetime of low-mass wind-driven supersoft X-ray binaries is $\lesssim 8$ million years. Wind-driven supersoft X-ray binaries with periods $\gtrsim 8$ hr cannot have steady-state hydrogen shell burning, but burn the accreted hydrogen in flashes. A system starting with a white dwarf $\gtrsim 1.25M_\odot$ may drive the white dwarf over the Chandrasekhar limit.

6. Discussion and conclusions

Observations show that in supersoft X-ray binaries the companion star is strongly irradiated by the accreting supersoft X-ray source, and that these systems have strong asymmetric mass outflow. Although there is no direct evidence that these winds are the result of the irradiation, we have shown that irradiation in supersoft X-ray binaries may lead to a strong stellar wind from the heated side of the secondary. Such a wind, assuming that it leaves the binary with specific angular momentum characteristic of the secondary, will drive mass transfer with a rate which is of the same order as the wind loss rate. For supersoft X-ray binaries with a low-mass companion ($M_2/M_1 < 2/3$), the resulting mass transfer is stable and with a rate sufficient to keep the binary in the stable (steady or recurrent) hydrogen burning regime.

The orbital periods of these low-mass supersoft X-ray binaries overlap those of the supersoft X-ray binaries with massive secondaries driven by thermal expansion of the donor star. It is therefore not possible to determine to which subclass a particular supersoft X-ray binary belongs just by inspection of its orbital period. An observational test would be the determination of the secondary's spectral type or a reliable determination of the binary's mass function. The thermal-timescale mass transfer requires a secondary with spectral type F or A, the wind-driven mass transfer spectral type M. Unfortunately, there is so far no unambiguous spectroscopic identification of the secondary in any supersoft X-ray binary. The mass functions of several supersoft X-ray binaries, determined using the He II $\lambda 4686$ emission line, are indeed very low (Crampton et al. 1987, 1996; Cowley et al. 1990; Beuermann et al. 1995). It seems unlikely that the He II $\lambda 4686$ line represents the true orbital motion of the compact accreting star (see Sect. 3.1), but if it is not too far off, a white dwarf primary would imply a low-mass expanded secondary. There is at least one known supersoft X-ray binary, 1E 0035.4-7230, which must have a low-mass M donor star because of its orbital period of 0.17 days.

Although wind-driven evolution is permitted for white dwarf binaries with low mass companions, it is clearly not the usual case: most systems in this region are ordinary cataclysmic variables. Some kind of ‘accident’ is required to trigger the high luminosities required for the wind-driven case. This is evidently a rare event, as shown by the fact that only one such low-mass system is so far identified with reasonable confidence. One possible way to enter the wind-driven evolution is a long phase of residual hydrogen burning after a mild shell flash, as proposed by Kahabka & Ergma (1997) for 1E 0035.4-7230. If such a flash lasts sufficiently long to change the mass transfer rate, the binary would unavoidably enter the wind-driven evolution regime. As discussed in Sect. 2, (low-)number statistics suggest that the supersoft X-ray phase of 1E 0035.4-7230 lasts longer than the $\sim 10^3$ yrs which are required to change the mass transfer rate significantly. Another possible way to obtain wind-driven mass transfer is a late helium shell flash of the cooling white dwarf, after the system has already come into contact as a cataclysmic variable. Such a late flash would cause the white dwarf to re-enter the horizontal shell-burning track and appear as a supersoft X-ray source.

Together with coalescing white dwarfs, supersoft X-ray binaries are the most promising progenitors of type Ia supernovae (e.g. Livio 1996). Thermal-timescale supersoft X-ray binaries, however, are only possible in fairly young stellar populations, while SNe Ia are also found in elliptical galaxies. Low-mass wind-driven supersoft X-ray binaries starting with an already massive white dwarf ($\gtrsim 1.25 M_\odot$) may drive the white dwarf over the Chandrasekhar limit and produce type Ia supernovae. Although the expected rate is difficult to estimate (depending on how often cataclysmic variables enter wind-driven evolution), these SN Ia progenitor candidates are also available for older populations, as in elliptical galaxies.

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