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Nova outbursts and the CV mass transfer rate spectrum

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Abstract. I investigate variations of the mass transfer rate in cataclysmic variables (CVs) induced by nova outbursts. The ballistic ejection of nova shells leads to a spectrum of observed transfer rates in a sample of systems with similar orbital period. I present an analytical model and show that the effect is negligible below the period gap and for CVs with high–mass white dwarfs. If the orbital period is longer than 3 hrs the mass transfer rate spectrum has an effective width in excess of an order of magnitude provided the mass accreted on the white dwarf prior to the runaway is in excess of a few $10^{-4}M_{\odot}$.

1. Introduction

A long–standing problem in the theory of cataclysmic variable (CV) evolution is the large scatter of observationally deduced mass transfer rates for systems with similar orbital periods (Warner 1995, Warner 1987, Patterson 1984). This appears to be in conflict with the well founded assumption that mass transfer in CVs is driven by orbital angular momentum losses such as gravitational radiation and magnetic braking.

Numerical calculations found that evolutionary tracks of CVs that form with different initial secondary masses converge rather quickly to an essentially unique track in e.g. orbital period P – mass transfer rate $\dot{M}$ space (e.g. Kolb & Ritter 1992). This track convergency arises as the angular momentum loss rate, and hence the transfer rate itself, depend on the current system parameters only. Low–mass main–sequence stars subject to mass loss rapidly settle at a new equilibrium radius R_2 (Stehle et al. 1996) which in turn is a function of the mass loss timescale. Therefore the orbital period $P \propto R_2^{3/2}$ of systems at the equilibrium radius is set by the mass transfer rate.

The track convergence implies that any variation of $\dot{M}$ at a fixed P should largely reflect the dependence of the angular momentum loss rate $\dot{J}$ on the white dwarf mass M_1. Since $-\dot{M}_2/M_2 \propto \dot{J}/J$, cf. eq (11) below, and $\dot{J}$ is independent of M_1 for mechanisms like magnetic braking which are rooted in the secondary star, the expected dependence is $\dot{M} \propto J^{-1} \propto \sqrt{M_1+M_2}/M_1$. This is too weak in view of the observed scatter of at least an order of magnitude.

2. Nova outbursts

Classical nova outbursts represent an obvious mechanism that could cause deviations of the transfer rate from the secular mean value set by the angular momentum loss rate. All CVs are expected to experience some form of nuclear burning of the accreted hydrogen–rich material on the surface of the WD. In most cases this should be explosive, in the form of a classical nova outburst, and hence accompanied by the ejection of the burning layers, the nova shell. The ejection occurs on a very short timecsale ($\lesssim$ years) and effectively leads to an instantaneous change of the orbital separation. The magnitude and sign of the change is determined by the amount of angular momentum the shell removes from the orbit. In most cases the separation increases, with a resulting drop of the mass transfer rate, as the photosphere of the secondary star suddenly finds itself a few scaleheights below the Roche lobe. The effect is largest if the ejection is ballistic, i.e. when additional angular momentum transfer due to dynamical friction of the secondary orbiting inside the nova envelope can be neglected.

Previous studies along these lines by MacDonald (1986), Livio, Govarie, & Ritter (1991), Schenker, Kolb, & Ritter (1998) dismissed the effect of nova outbursts on the mass transfer rate spectrum as unimportant.

Here I reconsider the maximum possible variation of the mass transfer rate due to novae. In Section 3 I present a simple analytical description. This is used in Section 4 to find the previously unrecognised optimum system parameter combination that maximises the spread of transfer rates in a sample of CVs with the same orbital period.

3. The analytic model

I shall make extensive use of standard expressions for the mass transfer rate and the Roche lobe index. The relevant notation is introduced in some detail in Sect. 3.1 and 3.2. In Sect. 3.3 I describe the application in the context of nova outbursts.

3.1. Stationary and instationary mass transfer

For convenience, let X be the magnitude of the mass transfer rate, i.e.

$$X = -\dot{M}_2 > 0, \tag{1}$$

where M_2 is the mass of the Roche–lobe filling secondary. In general, the transfer rate is a sensitive function of the difference between the stellar (photospheric) radius R and the Roche lobe radius R_L of the secondary. In the case of CVs a good approximation is

$$X = X_0 \exp\left\{\frac{R - R_L}{H}\right\}. \tag{2}$$

Here H is the photospheric pressure scale height of the secondary star, and $X_0 \simeq 10^{-8} M_\odot \mathrm{yr}^{-1}$ is the saturation value for optically thin mass transfer (e.g. Ritter 1988, Kolb & Ritter 1990).

Over periods that are short compared to the mass loss time $t_M = M_2/X$ the time derivative of the transfer rate (2) is

$$\dot{X} = X \, \frac{R}{H} \left\{ \frac{\dot{R}}{R} - \frac{\dot{R}_L}{R_L} \right\} \tag{3}$$

if X_0 and H are constant and $R - R_L \ll R$.

To make use of (3) we introduce the standard decomposition of radius derivatives into mass transfer–dependent and independent terms (e.g. Webbink 1985, Ritter 1996),

$$\frac{\dot{R}}{R} = \zeta \, \frac{\dot{M}_2}{M_2} + \left(\frac{\partial \ln R}{\partial t} \right)_{\rm th} \tag{4}$$

$$\frac{\dot{R}_L}{R_L} = \zeta_L \, \frac{\dot{M}_2}{M_2} + 2 \left(\frac{\dot{J}}{J} \right)_{\rm sys} . \tag{5}$$

Here the stellar mass–radius index ζ describes the adiabatic response of the donor star to mass loss, while $(\partial \ln R/\partial t)_{\rm th}$ denotes the relative radius change due to thermal relaxation. ζ_L is the Roche–lobe index and encompasses changes of the Roche lobe due to mass transfer and mass loss, and $\dot{J}_{\rm sys}$ are "systemic" orbital angular momentum losses unrelated to mass loss. J is the total orbital angular momentum.

With (4) and (5) equation (3) can be written as

$$\dot{X} = \frac{1}{\epsilon} \, \frac{X}{M_2} \left\{ \frac{M_2}{t_{\rm ev}} - DX \right\}, \tag{6}$$

where we have introduced the relative thickness of the photosphere of the donor star,

$$\epsilon = H/R, \tag{7}$$

(for CVs $\epsilon \simeq 10^{-4}$), the stability denominator

$$D = \zeta - \zeta_L, \tag{8}$$

usually of order unity, and the evolutionary timescale $t_{\rm ev}$,

$$\frac{1}{t_{\rm ev}} = \left(\frac{\partial \ln R}{\partial t} \right)_{\rm th} - 2 \left(\frac{\dot{J}}{J} \right)_{\rm sys} . \tag{9}$$

From (3) we see that mass transfer is stationary ($\dot{X} = 0$) if

$$\frac{\dot{R}}{R} = \frac{\dot{R}_L}{R_L}, \tag{10}$$

and (6) shows that in stationarity the mass transfer rate is

$$X_s = \frac{M_2/t_{\rm ev}}{D}. \tag{11}$$

Hence the transfer rate is of order $M_2/t_{\rm ev}$, unless the system is close to instability where D is small (see e.g. King & Kolb 1995 for a discussion of mass transfer instability).

3.2. The Roche lobe index

The stationary rate depends on the Roche lobe index. We obtain ζ_L by comparing the definition (5) with the time derivative of

$$J = M_1\, M_2 \sqrt{\frac{G\, a}{M}} \tag{12}$$

(a is the orbital separation, $M = M_1 + M_2$ the total binary mass), and by noting that R_L/a is a function of mass ratio $q = M_2/M_1$ only. The time derivative of the left hand side of (12) is the sum $\dot J_{\rm sys} + \dot J_{\rm ej}$ of systemic losses and losses due to mass that is lost from the system. We parametrize the latter by

$$\dot J_{\rm ej} = \nu\, \frac{J}{M}\, \dot M, \tag{13}$$

so that ν measures the specific angular momentum of the mass lost to the binary, in units of the mean specific angular momentum of the orbit. We further assume that the mass lost from the system is a fraction $\gamma > 0$ of the transferred mass,

$$\dot M = \gamma \dot M_2 \quad \text{or} \quad \dot M_1 = (\gamma - 1)\dot M_2. \tag{14}$$

Using the notation

$$\beta_2 = \frac{q}{R_L/a} \frac{{\rm d} R_L/a}{{\rm d} q} \tag{15}$$

this finally gives

$$\zeta_L \; = \gamma \left\{ (2\nu + 1)\frac{M_2}{M} + (-\beta_2 - 2)\frac{M_2}{M_1} \right\} + \zeta_c, \tag{16}$$

with ζ_c, the Roche lobe index for conservative mass transfer ($\gamma = 0$), given by

$$\zeta_c \; = 2\frac{M_2 - M_1}{M_1} + \frac{M}{M_1}\beta_2. \tag{17}$$

If $q \lesssim 1$ we have $R_L/a \simeq 0.46 (M_2/M)^{1/3}$ (Paczyński 1971), hence

$$\beta_2 \simeq \frac{1}{3}\, \frac{M_1}{M} \tag{18}$$

so that

$$\zeta_L \simeq \gamma \left\{ 2\nu \frac{M_2}{M} + \frac{2}{3}\frac{M_2}{M} - 2\frac{M_2}{M_1} \right\} + \zeta_c, \tag{19}$$

and ζ_c simplifies to the standard expression

$$\zeta_c \simeq 2\frac{M_2}{M_1} - \frac{5}{3}. \tag{20}$$

3.3. Nova outbursts

We now apply the above formalism to CVs that experience repeated nova outbursts. We consider both the actual evolution between outbursts, and the fiducial mean evolution. Quantities referring to this average over many nova outburst cycles will carry the index m.

We assume that the evolution between outbursts is conservative, i.e. $\zeta_L = \zeta_c$. The outburst itself is treated as an instantaneous ejection of the nova shell with mass $\Delta M_{\rm ej}$ and angular momentum $\nu J \Delta M_{\rm ej}/M$. If $\Delta M_{\rm ej}$ is related to the accreted mass $\Delta M_{\rm ign}$ needed to trigger the thermonuclear runaway as

$$\Delta M_{\rm ej} = \gamma \Delta M_{\rm ign} = -\gamma \Delta M_2 > 0 \tag{21}$$

($\gamma > 0$), then the Roche lobe index $\zeta_L = \zeta_m$ describing the mean evolution has just the form given by (16) or (19). In this fiducial mean evolution the sequence of discontinuous mass ejections is replaced by a continuous, istropic wind, such that $\dot{M}_{\rm wind} = \Delta M_{\rm ej}/t_{\rm rec}$, where $t_{\rm rec}$ denotes the outburst recurrence time. In the case where the ejected mass carries the specific orbital angular momentum of the white dwarf, $\nu = M_2/M_1$, (19) simplifies to

$$\zeta_m \simeq -\gamma \frac{4}{3}\frac{M_2}{M} + \zeta_c. \tag{22}$$

The corresponding long–term mean mass transfer rate ("isotropic wind–average") is

$$X_m = \frac{M_2/t_{\rm ev}}{D_m} \tag{23}$$

with $D_m = \zeta - \zeta_m$, cf. (11). The equivalent expression with $D_c = \zeta - \zeta_c$,

$$X_c = \frac{M_2/t_{\rm ev}}{D_c} \tag{24}$$

represents the conservative stationary rate the system would adopt in the absence of nova outbursts (if $D_c > 0$, i.e. if conservative evolution is dynamically stable).

Provided that the recurrence time is sufficiently short, D_c, $t_{\rm ev}$, ϵ and M_2 can be considered as constant between outbursts. Then (3) integrates to

$$X(t) = \frac{X_c}{1 + (X_c/X_0 - 1)\exp(-t/\epsilon t_{\rm ev})} \tag{25}$$

where X_0 is X at $t = 0$. This solution holds between $t = 0$, the time immediately after the outburst, and $t = t_{\rm rec}$, the time immediately before the next outburst. The behaviour of $X(t)$ depends on the sign of D_c. For $D_c > 0$ the transfer rate approaches the value X_c asymptotically, while for $D_c < 0$, $X(t)$ grows monotonically.

We note the equivalence

$$\int_0^{t_{\rm rec}} X(t) {\rm d}t = \Delta M_{\rm ign} = X_m \, t_{\rm rec}. \tag{26}$$

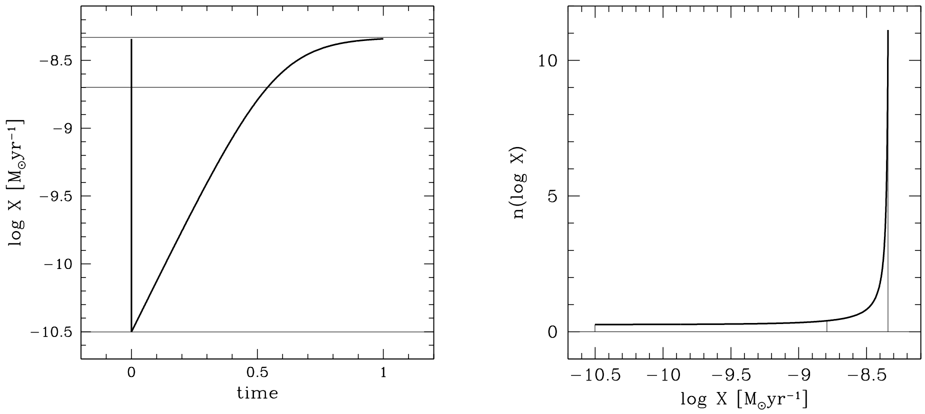

Figure 1. *Left panel:* Typical evolution of the mass transfer rate X between nova outbursts (here for $M_2 = 0.3 M_\odot$, $M_1 = 0.6 M_\odot$, $\Delta M_{\rm ign} = 3 \times 10^{-5} M_\odot$). The time is in units of the recurrence time. The upper solid line is the asymptotic value X_c, the lower solid line the time–average $X_m = 2 \times 10^{-9} M_\odot {\rm yr}^{-1}$. *Right panel:* The corresponding probability density distribution $n(\log X)$.

In what follows we assume that the thermal relaxation term $(\partial \ln R/\partial t)_{\rm th}$ which builds up in the donor star experiencing a discontinuous mass loss history due to nova outbursts is the same as in the time–averaged, continuous evolution. In other words, $t_{\rm ev}$ in (23) is the same as in (25). This has been confirmed by detailed integrations of the secular evolution using full stellar models.

The outburst amplitude X_f/X_0 of the mass transfer rate (X_f is X at $t = t_{\rm rec}$) can be calculated from the ejected envelope mass. From (2) we have

$$A = \frac{X_f}{X_0} = \frac{X\,(\text{pre-nova})}{X\,(\text{ post-nova})} = \exp\left\{ \frac{1}{\epsilon} \left(\frac{\Delta R_L}{R_L} \right)_{\rm out} \right\}, \tag{27}$$

where the index "out" refers to the nova outburst. Using (5) the total change $(\Delta R_L/R_L)_{\rm total}$ of the Roche lobe radius over one outburst cycle (outburst, followed by conservative evolution until the next outburst) can be written as

$$\left(\frac{\Delta R_L}{R_L} \right)_{\rm total} = \left(\frac{\Delta R_L}{R_L} \right)_{\rm out} + \left(\frac{\Delta R_L}{R_L} \right)_{\rm inter} = \zeta_m \frac{\Delta M_2}{M_2} + \left(\frac{\Delta J_{\rm sys}}{J} \right)_{\rm inter}, \tag{28}$$

while

$$\left(\frac{\Delta R_L}{R_L} \right)_{\rm inter} = \zeta_c \frac{\Delta M_2}{M_2} + \left(\frac{\Delta J_{\rm sys}}{J} \right)_{\rm inter}. \tag{29}$$

Here "inter" refers to the evolution between two outbursts, and $-\Delta M_2 = \Delta M_{\rm ign} > 0$ is the mass transferred in this phase. Hence we have

$$\left(\frac{\Delta R_L}{R_L} \right)_{\rm out} = (\zeta_m - \zeta_c) \frac{\Delta M_2}{M_2} = (D_c - D_m) \frac{\Delta M_2}{M_2}. \tag{30}$$

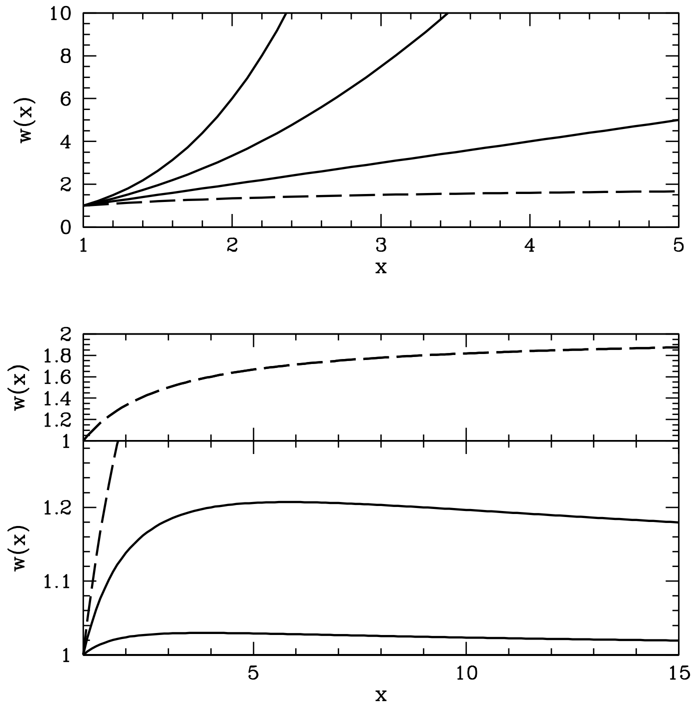

Figure 2. Half width $w(x)$ as a function of mass accumulation factor x, for different stability ratios α. The dashed curve is the limiting case $\alpha = 1$. Upper panel: $\alpha = -2, -1, 0, 1$ (from top to bottom); lower panel: $\alpha = 1, 1.5, 1.9$ (from top to bottom). See Sect. 4 for details.

Using (30) in (27) gives

$$A = \exp\left\{\frac{1}{\epsilon}\,(D_m - D_c)\,\frac{\Delta M_{\rm ign}}{M_2}\right\} \tag{31}$$

With (18) for β_2 and (22) for ζ_m this simplifies to

$$A = \exp\left(\frac{1}{\epsilon}\,\gamma\,\frac{4}{3}\,\frac{\Delta M_{\rm ign}}{M}\right) = \exp\left(\frac{1}{\epsilon}\,\frac{4}{3}\,\frac{\Delta M_{\rm ej}}{M}\right) \tag{32}$$

For this value of ν the Roche lobe of the secondary always increases as a result of a nova outburst, leading to a sudden drop of the transfer rate.

4. The mass transfer rate spectrum

Consider an ensemble of CVs with similar component masses (M_1, M_2) and mean mass transfer rate X_m at orbital period P. Then the actual (instantaneous) mass transfer rates in the ensemble are between X_0 and X_f. The probability density distribution $n(X)$ in X space is proportional to $1/\dot{X}$, so that the normalized

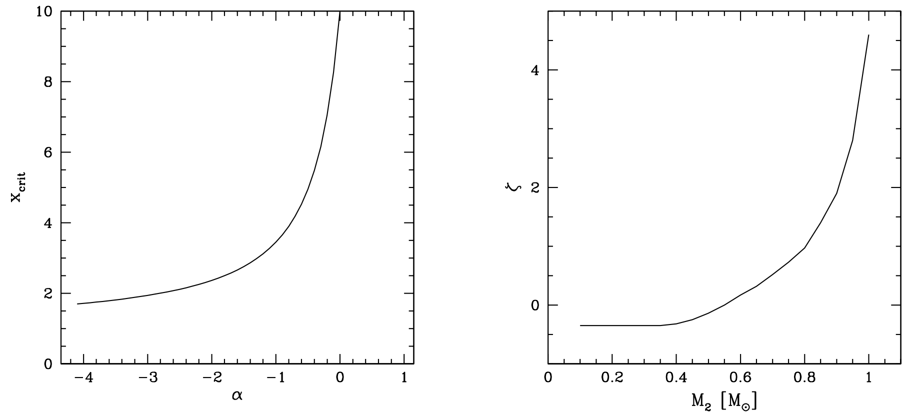

Figure 3. Left: $x_{\rm crit}(\alpha)$. Right: $\zeta(M_2)$ (see text for details).

mass transfer spectrum reads

$$n(X) = \frac{\mathrm{d}t}{\mathrm{d}X}(X)\,/t_{\rm rec} \tag{33}$$

($\int n(X)\mathrm{d}X = 1$). We use (24) in (6) to obtain for $\mathrm{d}t/\mathrm{d}X = 1/\dot{X}$

$$\frac{\mathrm{d}t}{\mathrm{d}X} = \epsilon\, t_{\rm ev}\,\frac{X_c}{(X_c - X)\;X}. \tag{34}$$

In integrated form, cf. (25), this is

$$t(X) = \epsilon\, t_{\rm ev}\,\ln\left\{\frac{X(X_c - X_0)}{X_0(X_c - X)}\right\}. \tag{35}$$

To calculate the mass transfer rate spectrum (35) or (34) for systems with specified masses and secular mean rate we need to find X_0. From (25) at $t = t_{\rm rec}$ we have

$$\frac{X_0}{X_c} = \frac{X_0/X_f - \exp\{-t_{\rm rec}/\epsilon t_{\rm ev}\}}{1 - \exp\{-t_{\rm rec}/\epsilon t_{\rm ev}\}}. \tag{36}$$

Using (31) to replace X_0/X_f and noting that

$$\frac{t_{\rm rec}}{\epsilon t_{\rm ev}} = \frac{\Delta M_{\rm ign}/X_m}{\epsilon t_{\rm ev}} = \frac{D_m}{\epsilon}\,\frac{\Delta M_{\rm ign}}{M_2}, \tag{37}$$

(cf. (26) and (23)) we finally find

$$\frac{X_0}{X_c} = \frac{\exp\{D_c\,\Delta M_{\rm ign}/\epsilon\, M_2\} - 1}{\exp\{D_m\,\Delta M_{\rm ign}/\epsilon M_2\} - 1} \tag{38}$$

In Fig. 1 we plot a typical mass transfer spectrum; note that the distribution shown is $\tilde{n}(\log X)$, i.e. the probability density distribution over $\log X$.

Equation (32) shows that the amplitude A can be very large if $\Delta M_{\rm ej}$ is large. But in many cases the systems spend most of their time close to the secular mean X_m, even if A is large. The corresponding mass transfer spectrum has a narrow spike at large transfer rates, and a long but hardly populated tail towards small X. The cases studied by Schenker et al. (1998) had an even narrower distribution than the example shown in Fig. 1.

To judge the effective width of the distribution in X-space we consider the half width w,

$$w \equiv \frac{X_f}{X_{1/2}}, \tag{39}$$

where $X_{1/2} = X(t_{\rm rec}/2)$ is the median of the mass transfer rate distribution (33), i.e.

$$\int_{X_0}^{X_{1/2}} n(X) \mathrm{d}X = \int_{X_{1/2}}^{X_f} n(X) \mathrm{d}X. \tag{40}$$

As a working criterion, we adopt the condition $w > 10$ to signal a "wide distribution". If we define the stability ratio

$$\alpha = \frac{2D_c}{D_m} \tag{41}$$

and the mass accumulation factor

$$x = \exp\left(\frac{D_m}{2\epsilon}\frac{\Delta M_{\rm ign}}{M_2}\right) \tag{42}$$

($x \geq 1$), then w can be written as

$$w(x) = \frac{x(1+x^{1-\alpha})}{1+x}. \tag{43}$$

For a given system, i.e. fixed mass ratio M_2/M_1, the stability factor α is fixed, while x depends on the ignition mass which we treat as a free parameter. Detailed model calculations of the thermonuclear runaway on the surface of hydrogen–accreting white dwarfs do not agree on this value. Observational estimates of nova ejecta tend to give much larger values than results from such calculations (see e.g. Starrfield 1999 and references therein). Hence we consider w as a function of x.

The shape of $w(x)$ depends on the value of the parameter α which describes the stability of mass transfer in the system. We require that CVs are dynamically stable with respect to the mean evolution, i.e. $D_m > 0$. With $\nu = M_2/M_1$ we also have $D_c < D_m$, hence $\alpha < 2$. The classical stability limit for conservative mass transfer ($D_c = 0$) is at $\alpha = 0$, while $\alpha < 0$ describes systems where mass transfer would be dynamically unstable in the absence of nova outbursts.

We plot $w(x)$ in Fig. 2. It is obvious that the character of $w(x)$ changes at $\alpha = 1$ (note that $w(1) = 1$ in all cases):

- If $2 > \alpha > 1$ the function $w(x)$ reaches a maximum at some $x = x_0$, and $w(x) \to 1$ for $x \to \infty$. We have $w \leq w(x_0) < 2$ in all cases. For the limiting case $\alpha = 1$, $w(x)$ is monotonically increasing and approaches the value $w = 2$ asymptotically. We conclude that for $\alpha \geq 1$ the mass transfer spectrum is always narrow.

- If $\alpha < 1$, $w(x)$ is monotonically increasing without bound. The slope increases with decreasing α. We have $w(x) = x$ for $\alpha = 0$. The condition $w > 10$ can be met in all cases, for any $x > x_{\rm crit}$, but $x_{\rm crit}$ is a function of α (Fig. 3, left panel). As α is in turn a function of mass ratio and adiabatic stellar index, the requirement $\alpha < 1$ is equivalent to the condition $q > q_{\rm crit}$, with

$$q_{\rm crit} \simeq 0.45 + 0.4 \left(\zeta_{\rm ad} + \frac{1}{3}\right). \tag{44}$$

5. Discussion

The physical meaning of the critical conditions ($q > q_{\rm crit}$, $x > x_{\rm crit}$) signalling "wide" mass transfer rate distributions becomes apparent when we translate them into constraints on system parameters. For a system with given secondary mass M_2 and donor mass M_1, the criterion $x > x_{\rm crit}$ is equivalent to a lower limit on the nova outburst trigger mass, $\Delta M_{\rm ign} > \Delta M_{\rm lim}$. In Fig. 4 we plot $\Delta M_{\rm lim}$ (in units of $\epsilon M_\odot$) as a function of M_2 for various M_1. In calculating these limits we assumed that $\gamma = 1$ ($\Delta M_{\rm ign} = \Delta M_{\rm ej}$), and that the secondary's adiabatic mass–radius index ζ is the same as for a ZAMS star with mass M_2 (shown in Fig. 3, right panel). The figure also gives a rough estimate of the orbital period the corresponding CV would have ($P/{\rm hr} \simeq 9(M_2/M_\odot)$, essentially independent of M_1).

The main conclusions from Fig. 4 are:

The nova–induced width of mass transfer spectra is insignificant below the CV period gap ($P \lesssim 2$ hr). These short–period systems are not close enough to instability, i.e. they do not conform to the mass ratio criterion $q \gtrsim 0.5$. For the same reason the nova–induced width is negligible whenever the white dwarf is massive ($M_1 \gtrsim 1 M_\odot$), whatever the orbital period.

Significant nova–induced widths can occur above the period gap, for systems with white dwarfs of mass $0.55 - 0.8 M_\odot$ and secondaries with mass close to the stability limit $M_2 \lesssim 0.7 M_1$. A necessary requiremment is, however, that nova trigger masses are in excess of a few times $10^{-4} M_\odot$. The actual value depends on the photospheric scaleheight H of the donor star, and increases with increasing white dwarf mass. Chiefly, significant widths can occur above the period gap, but should be least pronounced immediately above the gap and for long P.

A more detailed account of my analytical results, as well as the implementation of the scheme into a CV population synthesis code will be published in Kolb et al. (2000).

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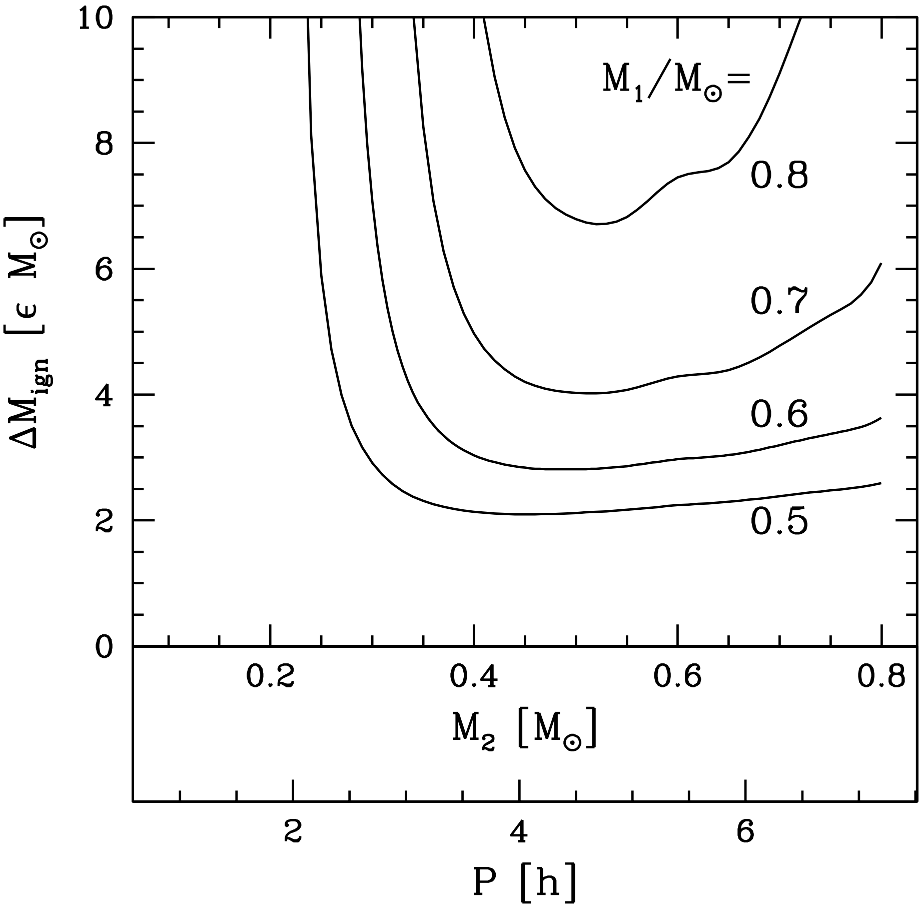

Figure 4. Lower limit on the nova trigger mass $\Delta M_{\rm ign}$ for nova–induced widths of the mass transfer rate spectrum in excess of a factor 10. $\Delta M_{\rm ign}$ is in units of $\epsilon \simeq 10^{-4}$ and shown as a function of donor mass M_2 or orbital period P, for different white dwarf masses M_1, as labelled.

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