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Classical Novae as Super-Eddington Steady States

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Abstract. The high luminosities and long decays of classical novae imply that they should be described as evolving super-Eddington (SED) steady states. We begin by describing how such states can exist—through the rise of a "porous layer" which reduces the effective opacity, and then discuss other characteristics of these states, in particular, that a continuum driven wind will arise. We then modify the stellar structure equations to describe these characteristics. The result is a modification of the classical core-mass—luminosity relation to include the super-Eddington state. The evolution of this state through mass loss describes classical nova light curves.

Key words. Stars: Novae

1. Introduction - Novae are SED steady states

One of the salient characteristics of classical novae is that the peak luminosity of most if not all eruptions is super-Eddington (SED). This can be seen in the peak luminosity of novae in M31 (see fig. 1).

Having SED luminosities is not contradictory to any classical notion, since SED luminosities can appear in dynamical systems (such as supernovae) which are far from a steady state. However, one would expect in such cases to develop bulk motions comparable to the local escape speed (since the effective gravity, including radiation, is similar to the actual gravity but reverse in direction). This implies that we can have a SED system but over a duration no longer than a time scale comparable to the typical size divided by the escape velocity. For the typical sizes of novae envelopes ($\sim 10^{12}$ cm) and typical wind velocities ($\sim$ 1000 km/s),

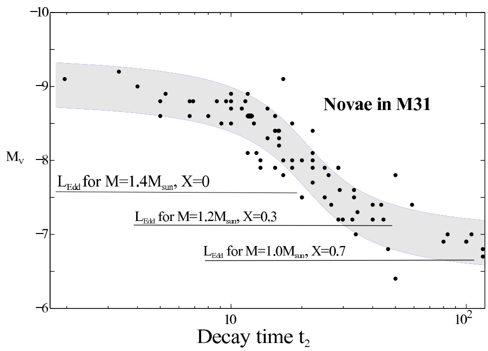

Fig. 1. The visible magnitude–decay rate relation of novae in M31, from della Valle & Livio (1995) The lines denote the Eddington luminosity limit for various cases. The upper line is the highest $L_{\rm Edd}$ theoretically possible for a Nova. The middle line is a more reasonable upper limit for typical fast novae.

one finds that novae could classically be SED for at most a few hours. And indeed, numerical simulations of novae can produce SED luminosities for a few hours but not longer (e.g. Prialnik & Kovetz 1992).

Nevertheless, a quick inspection of fig. 1 clearly reveals that the decay time is typically 10 days or longer. In fact, if one considers the bolometric luminosity, then the problem is aggravated since the decay rate of the total luminosity is typically a few times slower than the rate observed in the visible (Friedjung 1987; Schwarz et al. 2001; Shaviv 2001b). Clearly then, classical novae should be described as super-Eddington *steady states*.

This stands in stark contrast to the theoretical expectation. When the accreted material on a WD undergoes a thermonuclear runaway, it first passes a dynamic state during which the atmosphere is puffed up. In this stage, the luminosity can be easily SED. Because of that, a strong continuum driven wind should exist. However, once the system expands and stabilizes dynamically, it should follow the classical core-mass luminosity relation (Paczyński 1970). Since the CMLR is sub-Eddington, the only way to get a significant, optically thick mass loss is to drive a continuum driven wind on opacity maxima (e.g., Kato 1997). However, even in this case, the maximal mass loss possible is more than an order of magnitude smaller than many observed mass loss rates (Bath & Shaviv 1976; Shaviv 2001b).

The lower mass loss rate in the standard sub-Eddington picture has another interesting implication. It implies that the evolution is in many cases driven by the nuclear burning (i.e., changing of the envelope composition and not its mass (e.g., see Hernanz in this volume).

2. The Classical CMLR

Before we start modifying the standard CMLR, to see how it can describe super-Eddington states, we begin by revisiting the standard CMLR. This relation should describe any object with an inert degenerate core, and an envelope with a much smaller mass, and nuclear fusion at its base. Nova eruptions and post-AGB stars are the two classical examples of such objects. Empirically, one finds that such objects have an envelope mass independent luminosity for a wide range of envelope masses. Within this range, the outer radius of the object is much larger than the core radius. Below this $M_{\rm env}$ range, the CMLR breaks because the luminosity starts decreasing with the envelope mass, which extends to heights much smaller than the core radius. Above this range, the outer extent of the envelope is so large, that the escape velocity is comparable to the speed of sound, and the envelope evaporates. This is summarized in fig. 2.

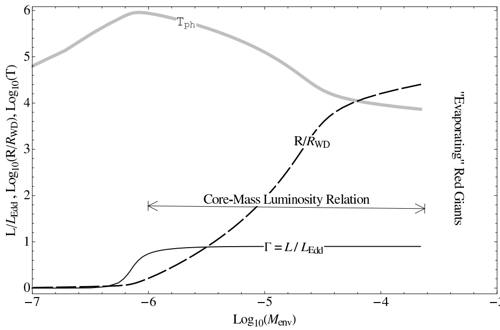

Fig. 2. The luminosity, photospheric temperature and photospheric radius as a function of $M_{\rm env}$ for the standard CMLR of a 1.2 $M_\odot$ object. For a finite range of envelope masses, the luminosity is almost constant. At larger envelope masses, the escape velocity is smaller than the speed of sound and the envelope evaporates. As lower envelope masses, the envelope "width" is smaller than the core radius, and the luminosity is smaller. The super-Eddington core mass (and envelope mass) luminosity relation is similar except that the "effective" luminosity is constant.

To understand what sets the luminosity and why it is not a function of the envelope mass, we should note that:

- To get an envelope which extends to large radii (much larger than the core radius), the sound speed at the base of the envelope has to be fixed, irrespective of the envelope mass. This gives a unique specific radiation entropy (s_r) luminosity relation through the dependence of the energy generation on the temperature.
- To get a consistent atmosphere which satisfies both the hydrostatic and radiative

transfer equations, there is another relation between s_r and $L/L_{\rm Edd}$, but one which is always sub-Eddington. Moreover, for a fixed opacity, s_r is the same everywhere in the envelope.

The $s_r - L$ relations have a unique intersection, which gives a specific L that only depends on the core mass. This is the origin of the CMLR.

3. How can SED states exist?

Before attempting to modify the CMLR, let us try to understand how SED states can exist. First, we should note that the opacity relevant to radiative transfer in *inhomogeneous* media is not the microscopic opacity. Instead, it is

$$\kappa_{\rm eff} = \frac{\langle F \kappa_\nu \rangle}{\langle F \rangle}, \quad (1)$$

where the average is a volume average and κ_ν is the opacity per unit volume (the extinction). Thus, any introduction of nonlinear structure will cause a change in the opacity. For a constant opacity (e.g., electron scattering) this always causes a reduction (Shaviv 1998).

The next point to consider is that as atmospheres approach $L_{\rm Edd}$ they become unstable, so that they naturally reduce their opacity. Such instabilities arise in Thomson scattering atmospheres (Shaviv 2001a) because of the vertical stratification. They can also arise on a larger range of length scales under more complex conditions, whether a non-constant opacity (Glatzel 1994), or a magnetic field (Arons 1992).

The above implies that if we wish to describe the global structure of luminous atmospheres, we should replace the microscopic opacity by an effective one. Since we do not now yet how the instabilities will saturate, we cannot know a priori the form for the opacity. In the following analysis, we will consider the simplest form which encapsulates the above physics: $\kappa_{\rm eff}(\Gamma) = \kappa_0/(1+\Gamma)$ where $\Gamma \equiv L/L_{\rm Edd}$. This gives no change in the effective opacity at small Eddington parameters Γ and a large enough reduction at high Γ's which keeps the system effectively sub-Eddington, that is, $\Gamma_{\rm eff} \equiv (\kappa_{\rm eff}/\kappa_0)\Gamma < 1$.

4. The modified CMLR

The next step is to consider the classical CMLR and modify it following the conclusions from the previous section. The first point to note is that the $s_r - L$ relation now expected form the radiative transfer and hydrostatic equilibrium will give the same $\Gamma_{\rm eff}$ (calculated with the effective opacity), however, because $\kappa_{\rm eff}$ is now reduced, the corresponding total luminosity is going to be higher, and it can be super-Eddington. This can be seen in fig. 3, where $L_{\rm eff} \equiv (\kappa_{\rm eff}/\kappa)L$ behaves the same as the luminosity in the classical CMLR (cf fig. 2), but the actual luminosity at the top of the atmosphere, $L_{\rm base}$, is much higher.

Although the internal structure (core + envelope) is very similar in the SED state, the appearance of SED states is going to be markably different from the sub-Eddington counterparts. The reason is that SED states necessarily accelerate a thick continuum driven wind (Shaviv 2001b). As a consequence, the photosphere and general appearance of the wind depends on the actual mass loss.

To derive the mass loss, we have to consider that the aforementioned opacity reduction can only take place as long as the nonlinear structure formed by the instabilities is optically thick. Once they become optically thin, the radiation cannot be funneled around the optical "obstacles". As a consequence, the anti-correlation between the flux and density needed for the reduction of the opacity disappears (see eq. 1).

Since the instabilities are expected to operate on a length scale comparable to the atmospheric vertical scale height, the effective opacity will approach the microscopic one where the optical width of a scale height becomes of order unity. From this, one can derive the expected mass loss (Shaviv 2001b):

$$\dot{m}_{\rm wind} = \mathcal{W} \frac{L - L_{\rm Edd}}{c v_s}, \quad (2)$$

where v_s is the sound speed at the sonic point. $\mathcal{W}$ is a constant (or weak function of Γ) which is expected to be of order unity. The mass loss derived from this mass-loss luminosity relation is consistent with η-Carinae and the winds of novae (Shaviv 2001b).

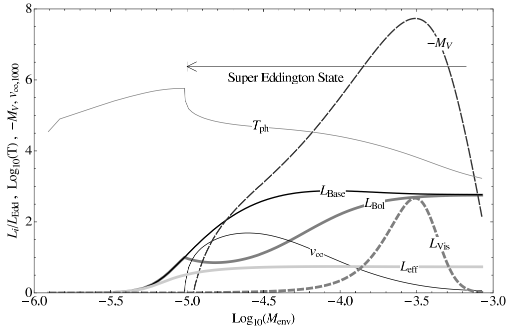

Fig. 3. Different characteristics of a shell burning object with a WD core of mass 0.8 $M_\odot$, as a function of the envelope mass, assuming $X = 0.3$ and $Z = 0.05$.

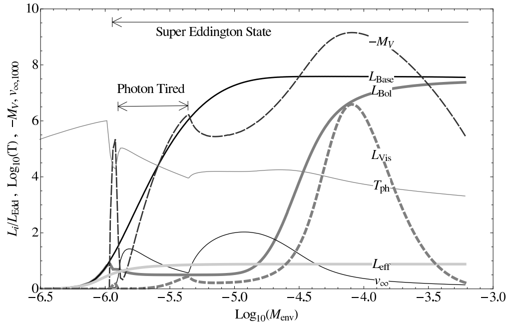

Fig. 4. The same as fig. 3, but for a 1.2 $M_\odot$ WD.

The resulting SED CMLR, which considers the modified opacity and the continuum driven wind is summarized, for two cases in figs. 3 and 4.

The steady states are obtained by solving the full stellar structure equations (though only with a Thomson scattering opacity). This is achieved by carrying out two iterations.

First, a radius for the sonic point is fixed while initial guesses for the luminosity and temperature at $r_{\rm sonic}$ are guessed. From the mass loss luminosity relation, one can derive the density and integrate outwards to infinity. The temperature at $r_{\rm sonic}$ is then iterated for until the photospheric black body conditions are met. This consistent wind solution is then integrated down to the WD radius. To obtain a consistent solution, the luminosity at $r_{\rm sonic}$ should be iterated for until $L(r_{\rm WD}) = 0$.

Note that we have two nested iterations (instead of iterating for T and L, or other two variables, simultaneously). It was found that the convergence of this algorithm is relatively fast and very stable.

Figs. 3 and 4 reveal several interesting characteristics of the SED states. First, as mentioned above, there are several relevant luminosities, not just one. The effective luminosity $L_{\rm eff}$ remains similar to the luminosity of the standard CMLR. However, the real luminosity $L_{\rm base}$ can be much higher. Moreover, because the wind is heavy, a large fraction of this luminosity can be used to accelerate the wind, which is why the observed $L_{\rm Bol}$ can be significantly smaller than $L_{\rm base}$.

Another interesting aspect can be seen in the $M = 1.2\ M_\odot$ case. For some envelope masses, the states become "photon tired". This condition arises because the energy required to carry the predicted mass loss out of the gravitational potential well is larger than the energy available at the base, and as a consequence, the wind stagnates (Owocki & Gayley 1997). The behavior of such systems was studied by van Marle et al. (2009). It was found that a new layer with a hierarchical set of shocks is formed. The layer effectively remains hydrostatic and it forms a new effective sonic point higher up in the potential well, were the mass loss would not be photon tired anymore. The layer does so because the supersonic material moving upwards advects the required energy, without the corresponding mass loss since there is a corresponding flow downwards.

5. Light curves from state evolution

Although the initial TNR of novae should be described dynamically, once the novae stabilize, they are described by steady states. However, these steady states do evolve in time. This evolution can translate into the nova light curve.

The evolution between different states can arise from the wind or the nuclear burning. The wind implies that as time progresses, the envelope mass simply decreases according to $\dot{m}_{\rm wind}$. Nuclear burning can cause two effects. The burning is responsible for a change in the

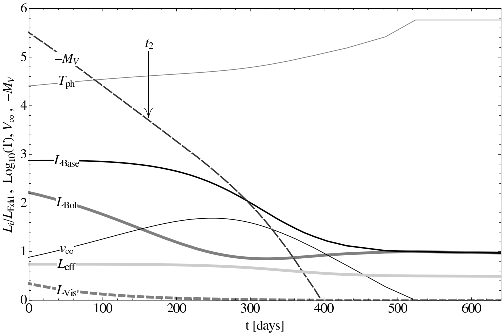

Fig. 5. A light curve obtained by relating states with different envelope masses (see fig. 3) through the continuum driven mass loss. The WD is assumed to have an $10^{-4} M_\odot$ envelope at $t = 0$.

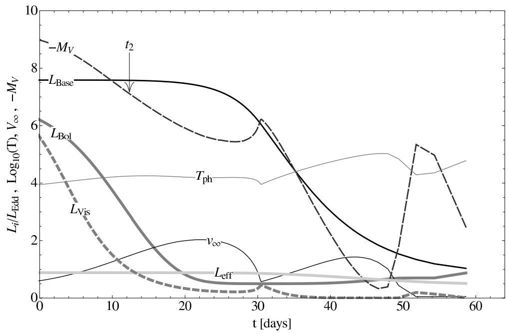

Fig. 6. The light curve obtained for the WD in fig. 4, assuming an initial envelope mass of $5 \times 10^{-5} M_\odot$.

composition. If the envelope is not convective, then the ashes will settle down on the WD and the envelope mass will decrease. If however the envelope is convective, the ashes will mix, and the net effect will be a reduction in the envelope's hydrogen fraction X.

We can define two time scales, one is due to the wind mass loss, $\tau_{\rm wind} \equiv M_{\rm env}/\dot{m}_{\rm wind}$, and a time scale to lose or change the mass through nuclear burning, $\tau_{\rm nuc} \equiv X M_{\rm env}/\dot{m}_{\rm nuc}$. Taking the above SED mass loss luminosity relation, and the nuclear energy production of $L = \epsilon \dot{m}_{\rm nuc} c^2$, we find

$$\frac{\tau_{\rm nuc}}{\tau_{\rm wind}} = \frac{\mathcal{W} X (\Gamma - 1) \epsilon c}{v_s \Gamma}. \quad (3)$$

This ratio is always much larger than unity by several orders of magnitude, implying that the state evolution is almost entirely due to the wind driven mass loss, at least, as long as a SED wind exists.

Armed with this knowledge, we can now relate the SED CMLR into a light curve through the wind mass loss.

Fig. 5 depicts the resulting light curve obtained for an $M = 0.8\ M_\odot$ nova having an initial envelope mass of $10^{-4}\ M_\odot$. These conditions give rise to a slow nova with a t_2 of about 160 days. Fig. 5 is similar, except that it depicts the light curve for an $M = 1.2\ M_\odot$ nova having an initial mass of $5 \times 10^{-5} M_\odot$. Here the nova is very fast, with t_2 of 12 days. It is also interesting that it exhibits a second peak. This is similar to the observed light curve of Nova V2491 Cyg.

6. Summary and discussion

We have seen that classical nova eruptions are SED and their decay is much longer than the dynamical time scale. Consequently, they should be described as super-Eddington steady states. Presently, the only known way how to explain the existence of these super-Eddington states is by having the effective opacity decrease as the Eddington luminosity is approached. Although at first it may seem as an ad hoc assumption, it is a natural consequence of the instabilities expected to take place when the radiation pressure becomes dominant.

The super-Eddington states which are then obtained resemble the classical CMLR. The main exception is that the SED states have a very large mass loss. These has many interesting implications. For example, the luminosity $L_{\rm Bol}$ observed at infinity is smaller than the luminosity at the base. As a consequence, $L_{\rm Bol}$ is a function of the envelope mass (unlike the classical relation).

The high mass loss rate also implies that the time dependence is a consequence of the decreasing envelope mass and not due to nuclear burning.

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References

Arons, J. 1992, ApJ, 388, 561
Bath, G. T. & Shaviv, G. 1976, MNRAS, 175, 305
della Valle, M. & Livio, M. 1995, ApJ, 452, 704
Friedjung, M. 1987, A&A, 179, 164
Glatzel, W. 1994, MNRAS, 271, 66
Kato, M. 1997, ApJS, 113, 121
Owocki, S. P. & Gayley, K. G. 1997, in Astronomical Society of the Pacific Conference Series, Vol. 120, Luminous Blue Variables: Massive Stars in Transition, ed. A. Nota & H. Lamers, 121
Paczyński, B. 1970, Acta Astronomica, 20, 47
Prialnik, D. & Kovetz, A. 1992, ApJ, 385, 665
Schwarz, G. J., Shore, S. N., Starrfield, S., et al. 2001, MNRAS, 320, 103
Shaviv, N. J. 1998, ApJ, 494, L193
Shaviv, N. J. 2001a, ApJ, 549, 1093
Shaviv, N. J. 2001b, MNRAS, 326, 126
van Marle, A. J., Owocki, S. P., & Shaviv, N. J. 2009, MNRAS, 394, 595

DISCUSSION

ODED REGEV: Is it possible that the explosion gives rise to non-spherically symmetric dynamics and then long term expansion, such that the Eddignton luminosity is not a problem?

NIR SHAVIV: The initial dynamic part of the explosion could indeed be asymmetric, which could facilitate having super-Eddington luminosities. The problem arises when considering time scales much longer than any relevant dynamic time scale, such that we expect gravity to equilibrate the envelope into a spherical hydrostatic solution. There is no way, that I know of, which could explain how such *hydrostatically equilibrated* states give rise to super-Eddington luminosities, other than by having a reduced opacity.

JEAN-PIERRE LASOTA: What determines the size of the blobs?

NIR SHAVIV: The size of the blobs is determined by the most unstable modes. This of course depends on the instability which operates, but the most general one which operates also in Thomson scattering atmospheres, requires the atmospheric stratification for the modes to grow. As a consequence, the least stable modes are of order the scale height of the atmosphere.

PIETER MEINTJE: Can one relate the dispersion relation of the instabilities to the observational behavior during outbursts?

NIR SHAVIV: The typical nonlinear structure has a typical length scale comparable to the atmospheric scale height below the sonic point. Since it is much smaller than the radius and because the wind acceleration might destroy it, I am doubtful that one can observe this structure directly. However, indirectly, the structure determines the mass loss rate—the wind constant $\mathcal{W}$ depends on the least stable mode, and it is expected to be of order unity for structure of order the scale height. Smaller unstable modes would give rise to a larger $\mathcal{W}$.