THE SUPERLUMINOUS SN DES13S2cmm AS A SIGNATURE OF A QUARK-NOVA IN A HE-HMXB SYSTEM

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ABSTRACT

We show that by appealing to a Quark-Nova (QN) (the explosive transition of a neutron star (NS) to a quark star (QS)) occurring in a helium-high-mass X-ray binary (He-HMXB) system we can account for the lightcurve of the first superluminous SN, DES13S2cmm, discovered by the Dark-energy Survey. The NS’s explosive conversion is triggered as a result of accretion during the He-HMXB’s second common envelope (CE) phase. The dense, relativistic, QN ejecta in turn energizes the extended He-rich CE in an inside-out shock heating process. We find an excellent fit (reduced $\chi^2$ of 1.09) to the bolometric light curve of SN DES13S2cmm including the late time emission, which we attribute to black hole accretion following the conversion of the QS to a black hole.

Key words: supernovae: general – supernovae: individual (SN DES13S2cmm)

1. INTRODUCTION

DES13S2cmm is the first confirmed superluminous SN (SLSN) from the Dark Energy Survey (Papadopoulos et al. 2015). Its bolometric light curve (LC) shows a slowly declining tail when compared to other SLSNe. It also shows an initial hump (the first four data points prior to the peak; see Figure 1). Papadopoulos et al. (2015) explored two possible power sources ($^{56}$Ni decay and magnetar spin-down power) but both models poorly fit the LC. In this paper we show that a Quark-Nova (QN) occurring in a helium-high-mass X-ray binary (He-HMXB), which experiences a second common envelope (CE) phase, can account for the LC features of SN DES13S2cmm including the tail. The paper is organized as follows: in Section 2 we give a brief overview of the QN model and the occurrence of QNe in binary systems. In Section 3 we show the results of applying the QN in a He-HMXB system to SN DES13S2cmm. We provide a brief discussion and some predictions in Section 4.

2. QN MODEL

The QN is the explosive transition of a neutron star (NS) to a quark star (QS) (Ouyed et al. 2002). The conversion energy combined with the core collapse of the parent NS results in the ejection of $M_{QN} \sim 10^{-3} M_\odot$ of neutron-rich material (Keränen et al. 2005; Ouyed & Leahy 2009; Niebergal et al. 2010). This relativistic ejecta has a kinetic energy, $E_{QS}$, exceeding $10^{52}$ erg and an average Lorentz factor $\Gamma_{QN} \sim 10$. The QN can occur following the explosion of a massive single star or in a binary system following accretion from the companion. The key constraint is for the NS to reach deconfinement densities in its core by becoming massive enough (e.g., via fall-back during the SN explosion or accretion from a companion or via spin-down (Staff et al. 2006)). We define $M_{NS,c}$ as the NS critical mass above which quark deconfinement occurs in the NS triggering the QN (see Ouyed et al. 2013b for a recent review). $M_{NS,c}$ varies from $\sim 1.6 M_\odot$ up to $2 M_\odot$ and higher for different equations of state (e.g., Staff et al. 2006). However, we assume $M_{NS,c} = 2 M_\odot$ to account for the heavy NS recently observed by Demorest et al. (2010); $\sim 2 M_\odot$ (or heavier) QSs can exist when one considers interacting quarks (e.g., Alford et al. 2007).

2.1. QNe in Single-star Systems: Dual-shock QNe (dsQNe)

A dsQN happens when the QN occurs days to weeks after the initial SN. The delay allows the QN ejecta to catch up to and collide with the SN ejecta after it has expanded to large radii (Leahy & Ouyed 2008; Ouyed et al. 2009). The SN provides the material at large radius and the QN re-energizes it, causing a re-brightening of the SN. For time delays of days or less, the radius of the SN ejecta is small enough that only a modest re-brightening results when the QN ejecta hits the SN ejecta. However, the neutron-rich QN ejecta lead to unique nuclear spallation signatures (Ouyed et al. 2011). For longer time delays, the radius and density of the SN ejecta are such that extreme re-brightening is observed which could explain some SLSNe (Ouyed et al. 2012; Ouyed & Leahy 2013; Kostka et al. 2014). For time delays exceeding many weeks, the SN ejecta will be too large and diffuse to experience any re-brightening. The dsQN model has been successfully applied to a number of superluminous and double-humped supernovae (Ouyed et al. 2013a; see http://www.quarknova.ca/LCGallery.html for a picture gallery of the fits).

2.2. QNe in Binary Systems

A QN will more often occur in tight binaries. To reach $M_{NS,c}$ and experience a QN event, the NS needs to accrete from a Roche Lobe (RL) overflowing companion (Ouyed et al. 2011c, 2011d; Ouyed & Staff 2013; see also Ouyed et al. 2015) or during a CE phase of the binary as described in this work. A QN in a binary provides:

(i) A means (the relativistic QN ejecta) to shock, compress, and heat the NS companion or the CE. In the case of a QN occurring inside a CE (which is considered in this paper), the resulting envelope thermal energy is given by $E_{CE,th} \sim \zeta_{sh}, E_{QN}$ which yields an initial shock temperature $T_{QN,sh} \sim \zeta_{sh} \times 10^9 K \times E_{QN,52} \times (\Gamma_{QE}/\Gamma_{CE})$ for a He-rich envelope of mass $M_{CE}$ and mean molecular weight $\mu_{CE} = 4/3; E_{QN,52} = 10^{52}$ erg while $\zeta_{sh}$ is the shock heating efficiency. The CE kinetic energy is $E_{CE,th} = E_{CE,gas} + E_{CE,rad}$, shared between the gas and the radiation. We find that even for extreme shock efficiency $\zeta_{sh} \sim 1$ and for high compression ratio ($\sim 4T_{QN} \sim 40$; see Ouyed & Staff 2013), the timescale required to burn He
to O (e.g., using burning rates from Huang & Yu 1998) is too large to compare to the adiabatic expansion timescale of the CE.

(ii) The spin-down power from the QN compact remnant (the QS), which provides an additional energy source. The QS is born with a magnetic field of the order of $10^{15}$ G due to color ferromagnetism inherent to quark matter during the transition (Iwazaki 2005). In the scenario considered here where the QS accretes during a CE phase, the high accretion rates would lead to QS spin-up instead of spin-down and the QS would gain mass until it becomes a BH without entering a spin-down phase. Thus in the picture presented here, spin-down power can be ignored.

(iii) The presence of a gravitational point mass (the QS) which may slow down and trap some of the ejecta and provide accretion power (e.g., Ouyed et al. 2011c, 2011d, 2014, 2015). The conversion of the QS to a BH during the CE phase provides an additional source of energy (the BH-accretion phase) which powers the LC in addition to energization by the QN shock.

### 2.2.1. The QN in a He-HMXB System

In previous papers we considered the detonation of a CO white dwarf by the QN (what we called a QN-Ia; Ouyed & Staff 2013; Ouyed et al. 2014, 2015). In a QN-Ia, the system experiences a runaway mass transfer leading to the formation of a CO-rich torus surrounding, and in the close vicinity of, the exploding NS. In this paper we consider a QN in a He-HMXB. Such a binary is expected to form when an O/B-HMXB binary evolves through a CE phase of the supergiant stellar component. If the binary survives this first CE event, the resultant system would contain a He-rich core (Hall & Tout 2014) with a NS in a relatively tight orbit. For a large mass ratio, a runaway mass transfer is expected once the He-rich core overflows its RL leading to the onset of a second CE phase. The NS accretes from two mass reservoirs (the first CE and the second CE). For a NS born with an initial mass of $\sim 1.6 M_\odot$, accretion of $\sim 0.4 M_\odot$ during the two CE phases is necessary to drive it above $M_{\text{NS,c}}$; accreting $\sim 0.2 M_\odot$ of material per CE phase is not unreasonable and should be enough to eject the first CE (Brown 1995; Armitage & Livio 2000). If the inspiralling NS accretes enough mass during the second CE phase, it should drive the NS above $M_{\text{NS,c}}$ to undergo a QN explosion in an extended CE; the CE expands radially outward as the NS inspirals toward the core. Following the QN explosion the QS continues to inspiral into the core while gaining mass. Below we show that if the QS turns into a BH, BH-accretion can power the slowly declining tail of DES13S2cmm. A combination of QN shock heating and BH-accretion provides the best fit to DES13S2cmm.

### 3. APPLICATION TO DES13S2CMM

The free parameters in our model are (see Table 1): the CE mass ($M_{\text{CE}}$; its initial (extended) radius $R_{\text{CE,0}}$ at the time of QN shock breakout; the CE expansion velocity $v_{\text{QN}}$ induced by the QN shock and QN shock heating per particle $(3/2)k_B T_{\text{QN,sh.}}$; and $\alpha_{\text{QN,sh.}}$, the shock propagation parameter (see the Appendix). There are three parameters for the BH-accretion model based on the prescription of Dexter & Kasen (2013) (see their Appendix A): $L_0$ (erg s$^{-1}$), $\Gamma_0$ and $\alpha$ which together define the injection power $L(t) = L_0(t/t_0)^\Gamma$ with $t_0 = \Gamma_0 t_d$ and $t_d$ the photon diffusion time in the CE; BH-accretion turns on at $t = t_0$ so that $L(t) = 0$ for $t < t_0$. Table 1 shows the model’s best-fit parameters for the BH-accretion model alone and the for the two-component QN+BH-accretion model. The best-fits reduced $\chi^2$ values are given in the last column of the table.

#### 3.1. BH-accretion Fit

This model does not have the QN explosion if the NS experiences a transition to a BH directly while the system is in the second CE phase. The NS could turn into a BH prior to or during merging with the CE core which forms an accretion disk around the BH to power the LC. The best-fit is obtained for $t_d = 13.7$ days, $\Gamma_0 = 0.1$ (i.e., $t_0 = 5.6$ days), $L_0 = 6 \times 10^{44}$ erg s$^{-1}$ and $\alpha = 0.7$. The fit is shown in Figure 1 alongside the observations from Papadopoulos et al. (2015). The resulting best-fit reduced $\chi^2$ value is $\sim 2.77$. The initial hump is not well fit with this model.

#### 3.2. QN+BH-accretion Fit

Adding the QN shock heating yields a significant improvement in the LC fit with a reduced $\chi^2_{\text{red}}$ of $\sim 1.09$ (see Table 1). The BH-accretion phase is delayed from the QN event by the time required for the QS to turn into a BH, merge with the core, and trigger accretion. This time delay is $t_d$ at which point the CE has extended to a radius $R_{\text{CE,0}} + v_{\text{QN}} t_d$; $R_{\text{CE,0}}$ is the CE radius at QN shock breakout (see the Appendix). The LC fit of DES13S2cmm is shown in Figure 2 alongside the observations from Papadopoulos et al. (2015).

The QN shock energizes the He CE (of mass $M_{\text{CE}} = 2 M_\odot$, and radius $R_{\text{CE,0}} = 1350 R_\odot$) and yields the initial bright and short-lived hump; the corresponding initial CE temperature $T_{\text{CE,0}}$ is calculated from the QN shock heating ($T_{\text{QN,sh.}} = 3.7 \times 10^6$ K; i.e., $\gamma_{\text{sh.}} \sim 10^{-2}$ for $M_{\text{CE}} = 2 M_\odot$) by including radiation energy density (see the Appendix). The photon diffusion timescale is $t_d = (2/3) \sqrt{M_{\text{CE}}/\ell_{\text{Th}}}/(0.3 \beta c v_{\text{QN}}) = 12.6$ days with $\beta = 13.8$ (Arnett 1982), $c$ is the speed of light, and $\ell_{\text{Th}}$ the Thompson cross-section.
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Table 1
Best-fit Parameters for the DES13S2cmm LC in Our Model

| He-rich (i.e., Second) CE | QN | BH Accretion | Fit |
|---|---|---|---|
| $M_{\text{CE}} (M_\odot)$ | $R_{\text{CE,0}}$ (R$_\odot$) | $v_{\text{QN}}$ (km s$^{-1}$) | $T_{\text{QN,sh.}}$ (K) | $\alpha_{\text{QN,sh.}}$ | $t_0$ (days) | $y_0$ | $L_0$ (erg s$^{-1}$) | $n$ | $\chi^2_{\text{red.}}$ |
| BH-accretion | ... | ... | ... | ... | ... | 13.7 | 0.10 | $6 \times 10^{44}$ | 0.7 | 2.77 |
| QN+BH-accretion | 2 | 1350 | 40,000 | $3.7 \times 10^6$ | 3/2 | ...$^a$ | 0.45$^b$ | $2.8 \times 10^{44}$ | 0.8 | 1.09 |

Notes.

$^a$ Here, $t_0$ is given by $t_0 = (2/3)^{1/3}M_{\text{CE}}/v_{\text{QN}}$ = 12.6 days.

$^b$ This corresponds to $t_0 = y_0 t_d \approx 5.6$ days. The kinetic energy of the CE/ejecta after the QN is $E_{\text{CE,K}} = (1/2)M_{\text{CE,QN}}^2 = 3.2 \times 10^{52}$ erg.

4. DISCUSSION AND CONCLUSION

In this paper we showed that the LC of SN DES13S2cmm is best fit with the QN+BH-accretion model. In our model, the QN occurs during the second CE phase of He-HMXB system followed by a BH-accretion phase after the QN turns into a BH in the core of the CE. The QN shock re-energizes the extended CE, explaining the initial hump, while BH-accretion power nicely fits the slowly declining tail. No nuclear burning is triggered by the QN shock which means that there should be little or no He-burning products in SN DES13S2cmm and similar explosions.

The lack of He-HMXB systems when compared to the observed large population of HMXBs with Be-type stars (e.g., Raguzaeva & Popov 2005 and references therein) has been used as an observational argument in favor of mergers during the dynamically unstable mass transfer phase of the progenitors (e.g., Linden et al. 2012; see also van den Heuvel 1976). However, even in extreme cases, the theoretically predicted fraction of He-HMXB surviving the CE phase is still higher than the observed number (Linden et al. 2012). If the NS is born massive, it will likely accrete enough matter during the first CE phase to reach $M_{\text{NS,c.}}$. The resulting QN could remove enough matter to unbind the system and bypass the production of He-HMXBs. We thus speculate that the QN may be partly responsible for the rarity of He-HMXBs.

If the NS does not accrete enough mass to go QN in the second CE phase, it may reach the center and form a Thorne–Żytkow object (Thorne & Żytkow 1977). However, continued accretion should trigger a QN leading to the same outcome, i.e., Thorne–Żytkow objects could be short-lived. On the other hand, extreme accretion rates could turn the NS directly into a BH.

The idea of a QN in binaries has proven successful in accounting for some features of SNe Ia (Ouyed et al. 2014, 2015) and Gamma Ray Bursters (Ouyed et al. 2011d) and has been able to account for the LC of DES13S2cmm as shown here. Its ability in fitting properties of unusual SNe (see http://www.quarknova.ca/LCGallery.html) suggest that QNe may be an integral part of binary evolution; the QN could lead to novel and interesting evolutionary paths. Nevertheless, as we have stated before, our model relies on the feasibility of the QN explosion which requires sophisticated simulations of the burning of a NS to a QS which are being pursued. Preliminary simulations with consistent treatment of nuclear and neutrino reactions, diffusion, and hydrodynamics show instabilities that could lead to a detonation (Niebergal et al. 2010; see also Herzog & Röpke 2011 and Albarracin Manrique & Lugones 2015). We have already argued that a “core-collapse” QN could also result from the collapse of the quark matter core (Ouyed et al. 2013b) which provides another avenue for the QN explosion.

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APPENDIX

Due to the outward diffusion of photons, the photosphere is moving inward in mass coordinates, slowly at first but faster as the density decreases in time. The ejecta interior to the photosphere we refer to as the core. We will assume that the thermal energy in the exposed mass in the photosphere (as the cooling front creeps inward) is promptly radiated. The interplay between uniform expansion and radiation diffusion defines the evolution of the photosphere as

$$R_{\text{phot}}(t) = R_{\text{CE}}(t) - D(t),$$

where $R_{\text{CE}}(t) = R_{\text{CE,0}} + v_{\text{QN}} t$ and $R_{\text{CE,0}}$ the CE envelope radius at QN shock breakout (which corresponds to $t = 0$ in our model). Here $D(t)$ is the diffusion length

$$D(t)^2 = D_0^2 + \frac{c}{n_{\text{CE}}(t) \sigma_{\text{Th.}}} t,$$

where $n_{\text{CE}}(t) = N_{\text{CE}}/V_{\text{CE}}(t)$ is the number density in the CE. The total number of particles in the CE is $N_{\text{CE}} = (M_{\text{CE}}/\mu_{\text{CE}} m_{\text{H}})$ while $V_{\text{CE}}(t) = (4\pi/3)R_{\text{CE}}(t)^3$ is the volume extended by the CE and $m_{\text{H}}$ the hydrogen atomic mass. We define $D_0$ as the initial diffusion length scale by setting

![Figure 2. The QN+BH-accretion model fit (solid line) to the light curve of SN DES13S2cmm.](image-url)
\[ n_{\text{CE,0}} \sigma_{\text{Th}} D_0 \simeq 1, \] where \( n_{\text{CE,0}} = N_{\text{CE}} / V_{\text{CE,0}} \) and the initial volume \( V_{\text{CE,0}} = (4\pi / 3) R_{\text{CE,0}}^3 \).

The initial QN shock heating per particle \((3/2)k_B T_{\text{QN,sh}}\) is a free parameter; \( k_B \) is the Boltzmann constant. The heat is redistributed between gas and radiation to get the post-shock CE temperature \( T_{\text{CE,0}} \). The relevant equation is

\[ \frac{3}{2} k_B n_{\text{CE,0}} T_{\text{CE,0}} + a_{\text{rad}} T_{\text{rad}}^4 = \frac{3}{2} k_B n_{\text{CE,0}} T_{\text{QN,sh}}^4, \]

with \( n_{\text{CE,0}} \) the number density (of electrons and ions) and \( a_{\text{rad}} \) the radiation constant.

The subsequent evolution of the CE core temperature after the CE is fully shocked is given by

\[ T_{\text{CE}}(t) = T_{\text{CE,0}} \left( R_{\text{CE,0}} / R_{\text{CE}}(t) \right)^2 \alpha_{\text{QN,sh}}. \]

To account for a non-uniform initial temperature, we introduce \( \alpha_{\text{QN,sh}} \), which parameterizes complex shock physics beyond the scope of this work. \( \alpha_{\text{QN,sh}} = 0 \) corresponds to the case of an adiabatic expansion with spatially uniform initial \( T_{\text{CE,0}} \), i.e., the internal energy includes only gas internal energy with \( \gamma = 5/3 \), so as time increases, \( T_{\text{CE}}(t) \propto R_{\text{CE}}(t)^{-2} \). \( \alpha_{\text{QN,sh}} \) also allows for the presence of radiation to be accounted, e.g., for spatially uniform \( T_{\text{CE,0}} \) and \( \gamma = 4/3 \), one uses \( \alpha_{\text{QN,sh}} = 1 \). \( \alpha_{\text{QN,sh}} > 0 \) can also correspond to a radially decreasing initial CE temperature so that \( T_{\text{core}} \) decreases more slowly than it would be for uniform \( T_{\text{CE,0}} \).

The corresponding luminosity is

\[ L_{\text{QN}}(t) = c_{V,\text{tot}}(t) \Delta T_{\text{core}}(t) n_{\text{CE}}(t) 4\pi R_{\text{phot}}(t)^2 \frac{d \Delta T}{dt}, \]

where the total specific heat is \( c_{V,\text{tot}}(t) = c_V \gamma + c_V \gamma_{\text{rad}} + \frac{3}{2} k_B + a_{\text{rad}} T_{\text{rad}}^3 / n_{\text{CE}}(t) \). Here, \( \Delta T_{\text{core}} \sim T_{\text{core}} \) since the photosphere cools promptly (i.e., cooling time is much less than the diffusion timescale). The inward photospheric velocity in mass coordinates is \( d \Delta T / dt \). When \( n_{\text{CE}} \) is low, the \( T^3 \) term dominates so that the temperature is much lower than in the pure gas model. As \( n_{\text{CE}} \) increases, the gas energy density becomes more important and \( T \) rises.

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