The range of Type Ia supernova (SN Ia) properties and new observations of these explosions at very early times (Nugent et al. 2011; Foley et al. 2012) point to a diversity of progenitor types. The long-standing view (see review by Hillebrandt & Niemeyer 2000) has been that SNe Ia result from a C/O white dwarf (WD) that has accreted enough material to compress the center to densities and temperatures so large that carbon fusion is ignited as an uncontrollable runaway that leads to an explosion. This requires the WD to nearly reach the Chandrasekhar mass ($M_{Ch}$), strongly constraining the binary evolution scenarios (Livio & Pringle 2011). The single-degenerate (SD) scenario has a main-sequence or slightly evolved donor, while the double-degenerate (DD) scenario has another C/O WD as the donor—likely through a merger (Iben & Tutukov 1984; Webbink 1984). These two channels may both occur, but a significant problem for the SD channel as a dominant mechanism is the lack of observed interaction between the SN ejecta and the hydrogen expected to be present in such a system.

Recent observations have shed light on possible progenitors via evidence of interaction between the SN ejecta and circumstellar material (CSM). Samples of SNe Ia show a bias toward blueshifted spectral features, such as Na absorption, which are likely outflows from the progenitor systems themselves (Sternberg et al. 2011). Previously reported SNe Ia showing evidence of CSM include SN 2002ic (Hamuy et al. 2003), SN 2005gj (Aldering et al. 2006), SN 2006X (Patat et al. 2007), SN 2007le (Simon et al. 2009), and PTF 11kx (Dilday et al. 2012). We note that there are alternative interpretations for 02ic as a Type Ic (Benetti et al. 2006) and 05gj as having a luminous blue variable progenitor (Trundle et al. 2008). Events such as SN 2006X and PTF 11kx show multiple blueshifted absorption features that may be interpreted as previously ejected material from successive recurrent nova outbursts in a symbiotic system such as RS Oph (Patat et al. 2007, 2011).

We present models of spherically symmetric recurrent nova shells interacting with circumstellar material (CSM) in a symbiotic system composed of a red giant (RG) expelling a wind and a white dwarf accreting from this material. Recurrent nova eruptions periodically eject material at high velocities ($\gtrsim 10^4$ km s$^{-1}$) into the RG wind profile, creating a decelerating shock wave as CSM is swept up. High CSM densities cause the shocked wind and ejecta to have very short cooling times of days to weeks. Thus, the late-time evolution of the shell is determined by momentum conservation instead of energy conservation. We compute and show evolutionary tracks of shell deceleration, as well as post-shock structure. After sweeping up all the RG wind, the shell coasts at a velocity $\sim 100$ km s$^{-1}$, depending on system parameters. These velocities are similar to those measured in blueshifted CSM from the symbiotic nova RS Oph, as well as a few Type Ia supernovae that show evidence of CSM, such as 2006X, 2007le, and PTF 11kx. Supernovae occurring in such systems may not show CSM interaction until the inner nova shell gets hit by the supernova ejecta, days to months after the explosion.

Key words: binaries: symbiotic – circumstellar matter – novae, cataclysmic variables – shock waves – supernovae: general

Online-only material: color figures
\( \dot{M} \approx 10^{-20} \text{km s}^{-1} \), and the ejecta velocity, \( \approx 3500 \text{ km s}^{-1} \) (Buil 2006; O’Brien et al. 2006).

In this paper, we examine the consequences of recurrent novae on the circumbinary environment. In Section 2, we show that the dense CSM in a symbiotic system leads to short radiative cooling times of the ejecta and shocked wind, making the kinematics of the nova shell determined by conservation of momentum. In Section 3, we derive the equations of motion for the spherically symmetric evolution of the nova shell and provide example realizations. The time-dependent post-shock structure of the decelerating ejecta and swept-up wind is discussed in Section 4, along with possible instabilities in the ejecta and their effect on the long-term evolution of nova shells. Finally, Section 5 discusses the implications of our model on potential supernovae in symbiotic systems.

2. KINEMATICS OF THE NOVA SHELL

An important difference between the evolution of nova ejecta in symbiotic systems and those in CV systems is the large amount of CSM that the ejecta will interact within a symbiotic system. The dense CSM around a symbiotic system decelerates the nova ejecta and decreases the radiative cooling time by many orders of magnitude from the \( > 10^4 \text{ yr} \) typical of nova shells interacting with the interstellar medium (ISM; Moore & Bildsten 2011). These short cooling times will cause momentum (rather than energy) conservation to determine the kinematics of the nova shell after a few weeks.

Although there is evidence that SyRNe such as RS Oph exhibit asymmetric mass ejection (O’Brien et al. 2006; Rupen et al. 2008), we consider a spherically symmetric model for simplicity. Between recurrent novae (with period \( t_{\text{rec}} \)) the RG donor is ejecting mass at a rate \( \dot{M}_w \) and velocity \( v_w \), creating a density profile

\[
\rho_e(r) = \frac{\dot{M}_w}{4\pi v_w r^2} = 2.2 \times 10^{-14} \text{ g cm}^{-3} \left( \frac{\dot{M}_w}{10^{-6} \text{ } M_\odot \text{ yr}^{-1}} \right) \times \left( \frac{v_w}{10 \text{ km s}^{-1}} \right)^{-1} \left( \frac{r}{1 \text{ AU}} \right)^{-2},
\]

out to a distance \( r_{\text{max}} = 6.3 \times 10^{14} \text{ cm} (v_w/10 \text{ km s}^{-1}) (t_{\text{rec}}/20 \text{ yr}) \) immediately before the next nova. This density profile will be perturbed by the presence of the accreting WD, as shown in the simulations by Walder et al. (2010), but remains nearly axisymmetric for slower wind velocities \( (v_w < 20 \text{ km s}^{-1}) \). Each nova ejects a mass \( M_{ej} \), which we scale as \( M_{ej} = f \dot{M}_w t_{\text{rec}} \), where \( f \) is a measure of both the accretion and explosion efficiency. We take \( f = 0.1 \) as our fiducial, although various simulations show effective accretion rates between 10% and 2% of \( \dot{M}_w \) (Walder et al. 2008, 2010), with \( \sim 90\% \) of the accreted material being ejected during a nova for RS Oph-like systems (Hachisu & Kato 2001).

Our model for the evolution of the nova ejecta is motivated by observations and simulations of outbursts in RS Oph. Novae with short recurrence times also have short mass-ejection timescales (Yaron et al. 2005), \( \sim 1-5 \text{ days} \), for recurrent novae such as RS Oph. Models of novae by Shen & Bildsten (2009) also show that short recurrence times require finely tuned mass-accretion rates. The mass-loss rate of the RG in RS Oph is \( \sim 10^{-6} \text{ } M_\odot \text{ yr}^{-1} \) (Rupen et al. 2008), which is consistent with a 20 year recurrence time given an accretion efficiency of \( f = 0.1 \). There are a range of measurements of the ejecta mass itself. The ejecta mass of the 1985 outburst was measured to be \( \sim 10^{-6} \text{ } M_\odot \) (O’Brien et al. 1992). Sokoloski et al. (2006) infer the ejecta mass of the 2006 outburst to be \( \sim 10^{-7} \text{ } M_\odot \) due to the quick onset of a Sedov–Taylor phase, while Vaytet et al. (2011) estimate it as \( (2-5) \times 10^{-7} \text{ } M_\odot \) from long-term simulations of the X-ray emission. Theoretical light curves of the 2006 nova by Hachisu et al. (2007) are best fit by an ejecta mass of \( (2-3) \times 10^{-6} \text{ } M_\odot \). The inferred ejecta mass in our model, neglecting the \( \sim 10\% \) of accreted material that may remain on the WD (Hachisu & Kato 2001), is \( M_{ej} = f \dot{M}_w t_{\text{rec}} \approx 2 \times 10^{-6} \text{ } M_\odot \), on the high side of estimates for RS Oph. Most measurements of the ejecta velocity are \( v_{ej} = 3000-3500 \text{ km s}^{-1} \) (Hjellming et al. 1986; Bul 2006), but models from Vaytet et al. (2011) argue that it is much higher, 6000–10,000 km s\(^{-1}\).

Early-time evolution of the ejection event is complex, requiring hydrodynamic wind–wind interactions (Vaytet et al. 2007, 2011), which we do not attempt to model. The 2006 outburst of RS Oph showed rapidly decelerating ejecta matching the self-similar Sedov–Taylor phase 3–10 days after the beginning of the outburst (Sokoloski et al. 2006; Bode et al. 2006), and quickly transitioning to the momentum-conserving phase after \( \sim 14 \text{ days} \) (Bode et al. 2006; Rupen et al. 2008). These observations indicate that the ejecta sweeps up enough mass in \( \sim 3 \text{ days} \) to get reverse-shocked and be in the self-similar phase. The cooling time of the post-shock material is thus \( \sim 14 \text{ days} \) in order to make the transition to a momentum-conserving phase.

The simplest model is to have the nova ejecta (here taken to be ejected all at once) coast into the \( \rho_e \) profile described above until it has sweep up mass equal to itself, at time \( t_{\text{sweep}} = 0.3 \text{ days} \) gives our fiducials of \( M_w = 10^{-6} \text{ } M_\odot \text{ yr}^{-1}, v_w = 10 \text{ km s}^{-1}, v_{ej} = 3000 \text{ km s}^{-1}, \) and \( t_{\text{rec}} = 20 \text{ yr} \). A slightly more realistic model is to remove the core of the wind profile so the ejecta starts encountering mass after it has traveled roughly the orbital separation, \( a \sim 0.5 \text{ AU} \). Doing so increases \( t_{\text{sweep}} \) to 3.0 days. After this time, the shell is in the self-similar Sedov–Taylor phase until radiative cooling makes the energy-conserving assumption invalid.

While in the Sedov–Taylor phase, the position and velocity of the forward shock are given by (Chevalier 1982)

\[
R_s = 9.8 \text{ AU} \left( \frac{f}{0.1} \right)^{1/3} \left( \frac{t_{\text{rec}}}{20 \text{ yr}} \right)^{1/3} \left( \frac{v_w}{10 \text{ km s}^{-1}} \right)^{1/3} \times \left( \frac{v_{ej}}{3000 \text{ km s}^{-1}} \right)^{2/3} \left( \frac{t}{5 \text{ days}} \right)^{2/3},
\]

\[
v_s = 2300 \text{ km s}^{-1} \left( \frac{f}{0.1} \right)^{1/3} \left( \frac{t_{\text{rec}}}{20 \text{ yr}} \right)^{1/3} \left( \frac{v_w}{10 \text{ km s}^{-1}} \right)^{1/3} \times \left( \frac{v_{ej}}{3000 \text{ km s}^{-1}} \right)^{2/3} \left( \frac{t}{5 \text{ days}} \right)^{-1/3}.
\]

The post-shock material is assumed to be fully reverse-shocked in this phase, so we do not follow any transient reverse shocks. The immediate post-shock particle density at time \( t \) is given by the strong shock jump conditions (Draine 2011)

\[
n_e = n_H = 9.1 \times 10^8 \text{ cm}^{-3} \left( \frac{v}{0.6} \right)^{-2/3} \times \left( \frac{M_w}{10^{-6} \text{ } M_\odot \text{ yr}^{-1}} \right) \left( \frac{v_w}{10 \text{ km s}^{-1}} \right)^{-5/3} \times \left( \frac{v_{ej}}{3000 \text{ km s}^{-1}} \right)^{-4/3} \left( \frac{t_{\text{rec}}}{20 \text{ yr}} \right)^{-2/3} \left( \frac{t}{5 \text{ days}} \right)^{-4/3}.
\]
where $\mu$ is the mean molecular weight of the material. The post-shock temperature is

$$T_s = 1.4 \times 10^7 \text{ K} \left(\frac{\mu}{0.6}\right) \left(\frac{v_s}{1000 \text{ km s}^{-1}}\right)^2.$$  

(5)

The post-shock cooling is roughly isobaric (Bertschinger 1986), so the cooling time of the post-shock material is

$$t_{cool} = \frac{n k T_s}{(\gamma - 1) \Lambda},$$  

(6)

where $\gamma = 5/3$ is the adiabatic index of the gas and $\Lambda$ is the cooling function. Using the approximation for the cooling function of $\Lambda/(n_e n_H) = 1.1 \times 10^{-22} T_s^{-0.7} \text{ erg cm}^3 \text{ s}^{-1}$ ($T_s$ is the temperature in units of $10^6$ K), valid for $10^5 K < T < 10^{7.3}$ K (Draine 2011), we calculate

$$t_{cool} = 36 \text{ days} \left(\frac{\mu}{0.6}\right)^{2.7} \left(\frac{f}{0.1}\right)^{1.8} \left(\frac{v_w}{100 \text{ km s}^{-1}}\right)^{2.8} \times \left(\frac{v_{ej}}{3000 \text{ km s}^{-1}}\right)^{3.6} \left(\frac{M_w}{10^{-6} M_\odot \text{ yr}^{-1}}\right)^{-1} \times \left(\frac{T_{rec}}{20 \text{ yr}}\right)^{1.8} \left(\frac{f}{t_{rec}}\right)^{0.2}.$$  

(7)

Using numerically computed cooling functions (Gnat & Sternberg 2007), rather than the power-law fit used above, reduces the cooling time to $t_{cool} \approx t$ for the first 16 days of evolution, after which $t_{cool} \ll t$, roughly agreeing with the cooling time inferred from shock deceleration measurements of RS Oph outlined above.

3. MOMENTUM-CONSERVING EVOLUTION

We now derive the equation of motion for the momentum-conserving phase. As will be shown in Section 4, the rapid cooling of the shocked material at late times causes most of the ejecta to be moving at the same velocity as the shock front, $v_s$. The initial momentum of the system is split between the momentum in the ejecta, $p_{ej} = f M_{w,rec} v_{ej}$, and that in the wind, $p_w = (1 - f) M_{w,rec} v_w$, where $f$ is the mass-sweep up the wind ejected since the previous nova, $t_{cool}$, and that in the wind is $p_w = (1 - f) M_{w,rec} v_w$. We immediately derive the final coasting velocity of the shell after it has swept up all the wind, using $p_{final} = M_{w,rec} v_{coast}$, and thus

$$v_{coast} = f v_{ej} + (1 - f) v_w.$$  

(8)

This implies coasting velocities of $\approx 100 \text{ km s}^{-1}$, intermediate to both the wind and nova velocities. At a time $t$ after the nova event the ejecta is sweeping up the wind, and the total wind mass swept up when the shell is at radius $R_s(t)$ is

$$M_{sweep}(t) = \int_{v_{wind}}^{R_s(t)} 4\pi r^2 \rho_w(r) dr = \frac{M_w}{v_w}(R_s(t) - v_w t).$$  

(9)

The integration must start at $v_{wind}$ because that is the outer radius of the wind that was ejected since the nova outburst. From this, we define the column (number) density of the shell as

$$N = \frac{M_{sweep}(t)}{4\pi \mu m_p R_s(t)^2}. $$  

(10)

The equation of motion for the shell arises from conservation of momentum:

$$(M_{sweep}(t) + f M_{w,rec}) v_{ej}(t) + [(1 - f) M_{w,rec} - M_{sweep}(t)] v_w = M_{w,rec} [f v_{ej} + (1 - f) v_w],$$  

(11)

and thus

$$\dot{R}_s = \frac{-(1 - f)(v_{ej}(t) - v_w)^2}{(1 - f)(R_s(t) - v_w t) + f v_{w,rec}}.$$  

(12)

We give examples of $R_s(t)$ and $v_{ej}(t)$ in Figure 1.

We can also obtain a simpler equation of motion in the limit of negligible wind velocity ($R_s \gg v_w t$),

$$R_s(t) = f v_w t_{rec} \left(\frac{2 v_{ej} t}{f v_w t_{rec}} - 1\right),$$  

(13)

defining an evolution timescale

$$t_{evol} = \frac{v_{w,rec}}{2 v_{ej}} = 4 \text{ days} \left(\frac{f}{0.1}\right) \left(\frac{v_w}{10 \text{ km s}^{-1}}\right) \left(\frac{t_{rec}}{20 \text{ yr}}\right) \times \left(\frac{v_{ej}}{1000 \text{ km s}^{-1}}\right)^{0.5}.$$  

(14)

Thus, for early times ($t \ll t_{evol}$) $R_s(t) \propto t$, while at late times $R_s(t) \propto t^{1/2}$. Coincidentally, $t_{evol}$ is shorter than the mass-ejection timescale of the nova, $t_{ej}$, as well as $t_{cool}$ in the Sedov–Taylor phase, and the momentum-conserving solution is not valid at these early times. Most of the observed shell evolution should therefore be in the $R_s(t) \propto t^{1/2}$ phase. High-resolution radio observations of the 2006 outburst of RS Oph tracked the deceleration of the shell, with Rupen et al. (2008) finding $v_{ej} \propto t^{-0.52}$ in the period 14–27 days after maximum—agreeing with the kinematics predicted by a momentum-conserving phase after a brief Sedov–Taylor phase. These calculations also yield the time required to completely sweep up the wind ejected since the previous nova, $t_{coast}$. From that time onward, the shell of material simply coasts outward at velocity $v_{coast}$. We show in the following section that the swept-up material is in a geometrically thin shell at $R_s$ and has a nearly uniform velocity throughout. In order to illustrate the resulting diversity in expected column densities,
of ejecta energy, for plane in Figure 2. Typical ejecta energies of SyRNe are
N
. POST-SHOCK STRUCTURE AND DENSITY ESTIMATES DURING DECELERATION

We have derived the expected column densities and coasting velocities from the simplest momentum-conserving considerations. The resulting values agree with those inferred in the few clear observations of circumstellar shells (Dilday et al. 2012; Patat 2011; Patat et al. 2007; Hamuy et al. 2003). However, we have not calculated the thickness, \(\Delta R\), of this shell, which sets the number density that is critical to photoionization and recombination calculations (Simon et al. 2009). Such estimates require
a consideration of the post-shock structure in the deceleration phase.

4.1. Cooling Gas

Calculation of post-shock structure of the nova blast wave is divided into two main parts: cooling and cooled gas. Immediately behind the shock front is the region of active radiative cooling. Since we have already established that the cooling time is short ($t_{\text{cool}} \ll R_s/(v_s)$, we calculate the fluid structure in this phase using the steady-state approximation. The hydrodynamic equations are written in Lagrangian coordinates, and we work in spherical coordinates so that the radial fluid velocity is represented by the single variable $u_r$. In the steady-state approximation we are left with the single independent variable $r$, indicating the radius from the explosion center. This yields three first-order differential equations:

$$\frac{\partial \rho}{\partial r} = \frac{-(\gamma - 1)\Lambda / \rho_r - 2 \rho^2 u^2_r / r}{u^2_r - \gamma P / \rho},$$

$$\frac{\partial \mu_r}{\partial r} = -\frac{u_r}{\rho} \left( \frac{\partial \rho}{\partial r} \right) - \frac{2 \rho u_r}{r},$$

$$\frac{\partial T}{\partial r} = -\frac{\left( \frac{\partial P}{\partial \rho} \right) T \left( \frac{\partial \rho}{\partial r} \right) - \rho u_r \left( \frac{\partial u_r}{\partial r} \right)}{\left( \frac{\partial P}{\partial T} \right) \rho}.$$  (18)

The post-shock conditions at $R_s$ come from the strong shock jump equations: $\rho_1 = (\gamma + 1)\rho_0 / (\gamma - 1)$, and $P_1 = (\gamma + 1)\mu m_p v^2_s / (\gamma + 1)$, and the ideal gas equation of state $P_1 = \rho_1 k T_1 / (\mu m_p)$. Here, we use non-equilibrium cooling functions from Gnat & Sternberg (2007) since cooling can be rapid enough to throw ion abundances out of equilibrium. We integrate the fluid equations from $R_s$ going in until the cooling rate drops off at $T_c = 10^4$ K and radiative cooling is no longer important. The cooling time is always much less than the age of the shell, so the shocked wind also follows the momentum-conserving evolution derived above. As the post-shock gas cools, it slows down relative to the shock front and increases in density.

We can estimate the density increase in the cold gas by noting that the cooling is roughly isobaric. The post-shock pressure is $P_1 \approx \rho v^2_s$, which must equal the pressure of the gas after it has cooled, $P_c = \rho k T_c / (\mu m_p)$ (c subscript indicates values in the cold gas). In terms of the thermal velocity of the cold gas, $v_{th} = \sqrt{2k T_c / (\mu m_p)} = 21 \text{ km s}^{-1}$, this gives

$$\frac{\rho_c}{\rho_0} \approx \left( \frac{v_s}{v_{th}} \right)^2 = 20 \left( \frac{v_s}{100 \text{ km s}^{-1}} \right)^2.$$  (19)

Note that the shock velocity is much slower than the ejecta velocity for all but the earliest evolutionary phases. A significant density increase in the cooled gas is also seen in Figure 5 of Vaytet et al. (2011), although the simulation shown is at a much higher ejecta velocity ($10^5 \text{ km s}^{-1}$) and earlier time (3 days), so we can only make a qualitative comparison. From our derived density contrast, we can also estimate the thickness of the cold gas via $4\pi R^2 \rho_c dR \approx M_{\text{sweep}}(t)$ and thus

$$\frac{\Delta R}{R_s} \approx \left( \frac{v_{th}}{v_s} \right)^2.$$  (20)

The cold gas is therefore a thin, dense shell as compared to the post-shock gas at $R_s$.

4.2. Cold Gas Evolution

Behind the cooling layer is the layer of cold, swept-up material. The evolution of this material depends on the secular evolution of the shock front, so the steady-state approximation breaks down. One method to obtain the structure in the swept-up matter is to assume that the deceleration of the shock front is quickly transmitted to all the gas, so that the internal pressure gradient is determined by the evolution of the shock front, $g \equiv d\mu_r / dt = -dP / dr$. By switching to mass coordinate behind the shock, $m$, we can write

$$\frac{dP}{dm} \approx \frac{g}{4\pi r^2}.$$  (21)

where $m$ is the mass outside of radius $r$, $m = \int_r^{R_s} 4\pi r^2 \rho(r) dr$, so $m / M$ is a mass coordinate measured from inside the shock (so that the shock front is at $m / M = 0$) and $g$ is negative since the shell is decelerating. The entropy of each mass element of wind gets frozen after radiative cooling has ceased. We find the entropy structure, $S(m)$, by using the steady-state cooling calculation described in the previous section. Finally, we assume that most of the mass is concentrated in a thin shell near the shock front ($r = R_s$) and calculate the pressure structure from hydrostatic balance,

$$\int_{P(0)}^{P(m)} dP = g \int_0^m dm / 4\pi r^2.$$  (22)

$$\Rightarrow P(m) = P(0) + \frac{g M}{4\pi R_s^2} \left( \frac{m}{M} \right).$$  (23)

Combining this with the input entropy profile and the adiabatic relations $\rho \propto (P / P_0)^{\gamma / 5}$ and $T \propto (P / P_0)^{\gamma / (\gamma - 1)}$ gives us the full post-shock structure (except for the velocity), shown in blue lines in Figure 5. We compare this calculation to that in Bertschinger (1986), which assumes a self-similar solution in the cold gas. That solution does not reproduce the post-cooling entropy of the shocked wind at early times because it uses a single power-law evolution in time, while the solution is transitioning from a radiative shock to a momentum-conserving one. The thin-shell approximation is more accurate for the regime where the outward-moving shock has not yet reached $R_s \propto t^{1/2}$.

4.3. Instabilities and Long-term Evolution

Chevalier & Imamura (1982) investigated instabilities of radiative shock waves in the ISM. Their analysis showed that for a cooling function $\Lambda \propto T^\alpha$, oscillatory cooling instabilities appear for $\alpha < 0.4$ (fundamental mode) and $\alpha < 0.8$ (first overtone). Recall that a fit to the equilibrium cooling function at relevant temperatures had $\Lambda \propto T^{-0.5}$, indicating an instability. We note, as do Chevalier & Imamura (1982), that non-equilibrium effects significantly alter the cooling function (Gnat & Sternberg 2007), so a true time-dependent calculation is necessary to determine the presence of instabilities in such radiative shocks. Such time-dependent calculations have been carried out for one-dimensional piston-driven radiative shocks (Innes et al. 1987; Gaetz et al. 1988), showing that shock velocities above $v_s \approx 150 \text{ km s}^{-1}$ are “overstable,” having oscillations in shock position relative to the piston location. Thus, steady-state shock calculations are not suitable for investigating the luminosities and spectra during the cooling phase of a radiative shock in our model.
As pointed out in Wood-Vasey & Sokoloski (2006), these early-time CSM interactions could be avoided if a previous nova swept out a cavity that has been refilled for less than $t_{\text{rec}}$. Assuming an ejecta velocity of $v_{\text{SN}} = 10^4 \text{ km s}^{-1}$ for the SN, then in the first $t$ days it could sweep up a wind mass that had been ejected for the last $t_{\text{SN}}$ years, where

$$t_{\text{SN}} = 5.5 \text{ yr} \left( \frac{v_{\text{SN}}}{10^4 \text{ km s}^{-1}} \right) \left( \frac{v_w}{10 \text{ km s}^{-1}} \right)^{-1} \left( \frac{t}{2 \text{ days}} \right).$$  \hspace{1cm} (24)

Thus, $t_{\text{SN}}$ is on the order of the recurrence times for very high mass WDs (Yaron et al. 2005; Shen & Bildsten 2009), so even supernovae with exceptionally early radio observations such as SN 2011Fe could hide significant CSM if it were in an SyRn system. Of course, there are many other lines of evidence that rule out such a progenitor for SN 2011Fe (Bloom et al. 2012; Chomiuk et al. 2012; Horesh et al. 2012; Margutti et al. 2012; Nugent et al. 2011; Li et al. 2011).

From Figure 4, the location of the nova shell at $t_{\text{rec}}$ gives an upper limit on the timescale of interaction with the innermost nova shell of

$$t_{\text{inner}} = 12 \text{ days} \left( \frac{R_s}{10^5 \text{ cm}} \right) \left( \frac{v_{\text{SN}}}{10^4 \text{ km s}^{-1}} \right)^{-1}. \hspace{1cm} (25)$$

This timescale can vary greatly due to the variability in $R_s$, and we note that it is consistent with the 22 day and 60 day brightenings observed in the light curve of SN 2002ic (Hamuy et al. 2003; Wood-Vasey & Sokoloski 2006).

6. CONCLUSIONS

The origins of SNe Ia continue to be debated, with some showing evidence of an SD progenitor (Dilday et al. 2012; Simon et al. 2009; Patat et al. 2007) and others with strong evidence of a DD progenitor (Bloom et al. 2012; Chomiuk et al. 2012; Horesh et al. 2012; Margutti et al. 2012; Nugent et al. 2011; Li et al. 2011). There is thus mounting evidence of multiple channels to a Type Ia.

Decelerating nova shells in SyRNe can have velocities consistent with the few SNe Ia with CSM detections. Nova shells are thin and have high-density contrasts ($\sim 100 \times$) compared to the ambient medium, which are important for calculating atomic populations. Additional work on simulating the ionization states of these shells and subsequent radiative transfer during an SN is necessary in order to make detailed comparisons to specific spectra. Supernovae in SyRN systems could be detected both via time-dependent absorption lines in previous shells and via rebrightening of the light curve as SN ejecta hits the shells.

We note some shortcomings of this model for interacting SNe Ia. First, novae (and RS Oph in particular) are known to have asymmetric ejecta (Hjellming et al. 1986; O’Brien et al. 2006; Rupen et al. 2008), which are not accounted for in our model. An asymmetric outburst (and/or wind profile) can change the ratio of momentum in the wind to that in the ejecta, and thus the coasting velocities. The general picture of a decelerating shell governed by momentum conservation remains, with details such as the density contrast and coasting velocity depending on orientation.

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