OPTICAL TO X-RAY SUPERNOVA LIGHT CURVES FOLLOWING SHOCK BREAKOUT THROUGH A THICK WIND

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ABSTRACT

Recent supernova (SN) observations have motivated renewed interest in SN shock breakouts from stars surrounded by thick winds. In such events the interaction with the wind powers the observed luminosity, and predictions include observable hard X-rays. Wind breakouts on timescales of a day or longer are currently the most probable for detection. Here, we study the signal that follows such events. We start from the breakout of the radiation-mediated shock, finding that the breakout temperature can vary significantly from one event to another (10⁴ to 5 × 10⁶ K) due to possible deviation from thermal equilibrium. In general, events with longer breakout pulse duration, t_{bo}, are softer. We follow the observed radiation through the evolution of the collisionless shock that forms after the breakout of the radiation-mediated shock. We restrict the study of the collisionless shock evolution to cases where the breakout itself is in thermal equilibrium, peaking in optical/UV. In these cases the post-breakout emission contains two spectral components—soft (optical/UV) and hard (X-rays and possibly soft γ-rays). Right after the breakout pulse X-rays are strongly suppressed, and they carry only a small fraction of the total luminosity. The hard component becomes harder, and its luminosity rises quickly afterward, gaining dominance at ~10–50 t_{bo}. The ratio of the peak optical/UV to the peak X-ray luminosity depends mostly on the breakout time. In early breakouts (t_{bo} ≤ 20 days for typical parameters) they are comparable, while in late breakouts (t_{bo} ≥ 80 days for typical parameters) the X-rays become dominant only after the total luminosity has dropped significantly. In terms of prospects for X-ray and soft gamma-ray detections, it is best to observe 100–500 days after explosions with breakout timescales between a week and a month.

Key words: stars: mass-loss -- stars: winds, outflows -- supernovae: general -- supernovae: individual (PTF 09uj, SN 2006gy) -- X-rays: general

Online-only material: color figures

1. INTRODUCTION

The breakout of a supernova (SN) shock through the stellar surface has been an active field of analytic and numerical research for several decades (e.g., Colgate 1974; Weaver 1976; Falk 1978; Klein & Chevalier 1978; Imshennik et al. 1981; Enssner & Burrows 1992; Matzner & McKee 1999). Recent advancement of observational facilities led to the discovery of several shock breakout candidates (e.g., Campana et al. 2006; Soderberg et al. 2008; Gezari et al. 2008, 2010; Schawinski et al. 2008; Modjaz et al. 2009; Ofek et al. 2010; Arcavi et al. 2011), which motivated a revisiting of the theory of shock breakout (e.g., Chevalier & Fransson 2008; Katz et al. 2010, 2012; Pirz et al. 2010; Nakar & Sari 2010, 2012; Murase et al. 2011; Rabinak & Waxman 2011; Couch et al. 2011; Chevalier & Irwin 2011; Balberg & Loeb 2011; Moriya & Tominaga 2012).

The main observational challenge posed by shock breakouts from stellar surfaces is their short duration (no more than an hour even in the case of a red supergiant). However, in some cases massive mass loss prior to the SN explosion can extend the breakout signal, facilitating its detection. Ofek et al. (2010) and Chevalier & Irwin (2011) show that a breakout through a thick wind may occasionally extend long enough to account for the complete SN luminosity light curve and suggest that at least some Type IIn SNe are in fact powered by such breakouts. When a star goes through an SN explosion, a radiation-mediated shock (RMS) is driven through its envelope, accelerating through the sharp density drop near the stellar edge (Sakurai 1960). If the star is not surrounded by a thick wind, then the shock breaks out of the stellar surface, producing an intense short pulse of UV–X-ray photons. If, on the other hand, the star is surrounded by a wind with an optical depth between the stellar edge and the observer >c/v, where c is the speed of light and v is the shock speed, then the shock continues into the wind without releasing photons to the observer. At this point the shock starts decelerating and a reverse–forward shock structure is formed. The reverse shock is driven into the SN ejecta, which is characterized by a sharp rise in the amount of energy that is carried by slower moving material. The reverse shock is driving an increasing amount of energy into the shocked region, which keeps accumulating until the wind optical depth drops to c/v. Beyond this point the wind cannot contain anymore the radiation-mediated forward shock and the shock is breaking out of the wind, releasing all its energy as an intense pulse. This pulse can contain much more energy than a windless breakout, over timescales that are much longer. The general properties of a wind breakout were recently discussed by Ofek et al. (2010) and Chevalier & Irwin (2011), who find the pulse luminosity and temperature assuming thermal equilibrium, and by Balberg & Loeb (2011), who also show how the progenitor mass and explosion energy can be constrained by the pulse observational properties.

If the wind density does not fall abruptly, the energy behind the shock keeps growing after breakout, and the characteristic evolution of the resultant signal supplies an important observational probe for the interaction of the forward shock with the wind. Katz et al. (2011) showed that if the breakout of the RMS occurs within the wind, i.e., the matter comprising the breakout...
layer is spread over a scale comparable to the breakout radius rather than being concentrated within a thin layer, then a collisionless shock is bound to develop and replace the RMS in accelerating the matter ahead. This collisionless shock heats the immediate post-shock electrons to $k_B T_e \sim 60$ keV (Katz et al. 2011; Murase et al. 2011), where $T_e$ is the electron temperature and $k_B$ is the Boltzmann constant. The cooling processes of these electrons affect the observed temperature (and possibly the luminosity evolution), which is further influenced by the interplay between photons and electrons along the photons’ diffusion path to the observer.

Here, we examine the evolution of the luminosity and the observed temperature at and following a wind shock breakout. We consider only breakouts through a wind that is dense enough to produce a breakout pulse lasting days or longer, which are easier to detect. We restrict the discussion on the observed radiation during the collisionless shock evolution to scenarios where the unshocked electrons ahead of the shock and the radiation that diffuses through them on its way to the observer are in thermal equilibrium. This condition implies shock velocities $\lesssim 10^5$ km s$^{-1}$, where the exact limiting velocity can vary with the gas metallicity. Finally, we consider only cases where the shocks are cooling fast during the entire evolution, i.e., the shock-heated plasma cools within a dynamical time. This condition is satisfied for regular SN parameters. We pay attention to hydrodynamic, diffusion, and cooling timescales involved and show that the breakout timescale is the most dominant parameter in determining the post-breakout evolution. We map the different types of post-breakout behaviors and describe their light curves and spectra evolution.

As this work was near completion, Chevalier & Irwin (2012) submitted a paper presenting some of the conclusions derived below in the context of explaining the low X-ray luminosity of SN 2006gy, observed by Chandra at an age of 3–4 months. The low luminosity is explained as the result of three factors: two, which are discussed in detail below, are the domination of inverse Compton (IC) over free–free emission in cooling the hot post-shock electrons and the energy degradation of the hard free–free emitted photons over the diffusion path. The third is photoabsorption of soft X-rays by partially ionized pre-shock gas. The latter process is not considered in this work.

2. HYDRODYNAMIC EVOLUTION AND CHARACTERISTIC SCALES

We consider the interaction of SN ejecta with a dense stellar wind that follows a standard profile $\rho_w = Dr^{-2}$, where $D$ is a constant and $r$ is the radius. This wind profile implies $\tau \propto r^{-1}$, where $\tau$ is the optical depth of the wind at the location of the shock. It also implies a photon diffusion timescale to the observer, which is independent of the radius (up to a logarithmic factor), $t_{\text{diff}} \sim \kappa D/c$, where $\kappa$ is the cross section per unit of mass (throughout the paper we assume that the gas that dominates the optical depth is fully ionized and take $\kappa = 0.34 \text{ cm}^2 \text{ g}^{-1}$).

During the interaction, a forward shock is driven into the wind while a reverse shock is driven into the ejecta. Chevalier (1982) derives a self-similar solution of the forward–reverse shock structure, assuming that both the ejecta and the wind density profiles are power laws:

$$r(t) \propto t^{(n-3)/(n-2)},$$

for $\rho_w \propto r^{-2}$ and $\rho_{\text{ej}} \propto r^{-n}$. The ejecta density profile through which the reverse shock is propagating is determined by the SN shock that crosses the stellar envelope. The density profile of the fastest moving ejecta can be approximated by a power law with an index $n \approx 10–12$ for all progenitor types. The ejecta density profile flattens toward the slower layers, which contain most of the ejecta mass and energy (e.g., Matzner & McKee 1999). As the reverse shock approaches these layers, the evolution is no longer strictly self-similar, but it can be approximated by the self-similar solution using $n = 7$ (e.g., Chevalier & Irwin 2011, 2012). Given that during the reverse–forward shock interaction phase $n \gtrsim 7$, we hereafter approximate $(n-3)/(n-2) \approx 1$ during this phase. Finally, once the reverse shock ends crossing the ejecta, only the forward shock remains, entering a Sedov–Taylor phase if the shock is adiabatic or a snowplow phase if it is radiative.

The density power-law indices of the ejecta and the wind are set by the pre-explosion stellar evolution and are not expected to vary much between various SN progenitors. The difference in the hydrodynamical evolution during the forward–reverse shock phase between various SNe is set by the two normalization factors of these density profiles. These could be determined by two observables such as the duration of the breakout pulse, $t_{\text{bo}}$, which is set by the wind density only, $t_{\text{bo}} \sim \kappa D/c$ (Ofek et al. 2010; Chevalier & Irwin 2011; Balberg & Loeb 2011), and the breakout luminosity $L_{\text{bo}}$. Below we present the evolution during the interaction phase using $t_{\text{bo}}$ and the shock breakout velocity $v_{\text{bo}}$ (instead of $L_{\text{bo}}$), which is related to the observables via $v_{\text{bo}} = ((\kappa/4\pi c)L_{\text{bo}}t_{\text{bo}}^{-2})^{1/3}$ (Balberg & Loeb 2011). The evolution after the interaction phase ends depends on three parameters, and we present it using $t_{\text{bo}}$, the total explosion energy, $E$, and ejecta mass, $M_{\text{ej}}$.

The breakout from the wind takes place once $t_{\text{bo}} \approx c/v_{\text{bo}}$, where $t_{\text{bo}}$ is the wind optical depth to the observer at the breakout radius. This sets the mass swept by the shock at the time of the breakout:

$$m_{\text{bo}} \approx 4\pi c/\kappa v_{\text{bo}}t_{\text{bo}}^2 \approx 5 \times 10^{-3} M_{\odot} v_{\text{bo}}^2 t_{\text{bo}}^{-2},$$

where $v_{\text{bo}} = v_{\text{bo}}/10^9$ cm s$^{-1}$ and $t_{\text{bo}},d = t_{\text{bo}}/1$ day. During the ejecta–wind interaction, the wind mass collected by the forward shock is comparable to the ejecta mass collected by the reverse shock. Therefore, $v_{\text{bo}}$ can be found using the velocity profile of an ejecta released by an SN explosion (e.g., from Nakar & Sari 2010, hereafter NS10), and requiring that a wind mass collected by the forward shock until the breakout, $m_{\text{bo}}$, is accelerated to $v_{\text{bo}}$ by the explosion. Equations (A2) and (A4) from NS10 provide the ejecta velocity profile $v(m) \approx 3 \times 10^8$ cm $s^{-1}(m/M_{\text{ej}})^{0.19/\beta-4/3}(E_{S1}/M_{\text{ej}})^{1/2}$, where $M_{10} = M_{\text{ej}}/10 M_{\odot}$, $E_{S1} = E/10^{51}$ erg, and $n$ (not to be confused with $n$) is the power-law index of the pre-explosion stellar density profile near the edge ($\bar{n} = 3$ for a fully radiative envelope and $\bar{n} = 1.5$ for a fully convective one). Plugging Equation (2) into this velocity profile, we find

$$v_{\text{bo}} \approx 10^9 \text{ cm s}^{-1} M_{\text{ej}}^{0.35} E_{S1}^{0.45} t_{\text{bo}}^{-0.2},$$

where the dependence on $\bar{n}$ is weak. This equation is applicable until $m_{\text{bo}} = M_{\text{ej}}$. It implies that for typical SNe (e.g., not

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3 The mass profile for $m \ll M_{\text{ej}}$ is $m \propto v^{-6+\delta_3}/10^{19}$, where $\bar{n}$ (not to be confused with $n$) is the power-law index of the pre-explosion stellar density profile near the edge. This profile translates, under free expansion, to $\rho_{\text{ej}} \propto t^{-3} v^{-3-(\delta_3+1)/10^{19}}$, or $n \approx 10$ for a radiative envelope and $n \approx 12$ for a convective one.
3. TEMPERATURE AND LUMINOSITY EVOLUTION

3.1. Bolometric Luminosity

The bolometric luminosity of a fast-cooling shock is $L \approx 2\pi r^2 \rho v^3 \propto v^3 \propto \alpha^{-3/(n-2)}$. Therefore, if the breakout takes place before $t_{SP}$, then

$$L(t) \sim \begin{cases} 5 \times 10^{51} t_{bo,d} \rho_{bo,9} v_{bo,9}^3 \left( \frac{t}{t_{bo}} \right)^{-0.3} \text{erg s}^{-1} & t_{bo} < t \ll t_{SP} \\ 2 \times 10^{52} t_{bo,d} M_{10}^{1/2} E_{51}^{3/2} \left( \frac{t}{t_{SP}} \right)^{-1.5} \text{erg s}^{-1} & t_{SP} \lesssim t, \end{cases}$$

while if $t_{bo} > t_{SP}$, then

$$L(t) \sim 5 \times 10^{53} t_{bo,d} \rho_{bo,9} v_{bo,9}^3 \left( \frac{t}{t_{bo}} \right)^{-1.5} \text{erg s}^{-1}; \quad t_{SP} < t_{bo} < t. \quad (7)$$

Note that once the radiation escape time becomes shorter than the dynamical time (i.e., after breakout), then as long as the shock is cooling fast, the bolometric luminosity is simply the kinetic energy flux through the shock. In contrast to regular SNe, it is independent of the gas opacity and thus insensitive to processes such as recombination or to the gas metallicity.

3.2. Breakout Temperature

The first signal to be observed is the breakout pulse, releasing over a diffusion time $t_{diff} \sim t_{bo}$ the energy accumulated during the interaction with the wind prior to the breakout. The pre-breakout forward shock is radiation mediated, but the photons that mediate the shock and dominate the energy density behind the shock cannot escape before breakout. As the breakout occurs, these photons start leaking toward the observer and their observed temperature is set by their interaction with the unshocked wind through which they diffuse.

The unshocked wind is heated by the diffusing radiation. It produces photons, mostly by free–free emission, which share energy with the diffusing radiation. The breakout temperature is therefore set by the ability of the unshocked wind to produce photons over the radiation diffusion time. The radiation energy density at any point along the diffusion path is $\epsilon_{rad} \sim (L \tau/c^2)$. We define $T_{\text{FB}} \equiv (\epsilon_{rad}/a)^{1/4}$, where $a$ is the radiation constant.

If photons are abundant enough behind the shock and the radiation is in thermal equilibrium, then the photon temperature at the shock is $T_{\text{FB}}$. The observer is seeing, however, a lower temperature. The reason is that thermal equilibrium is also kept at radii that are larger than the shock radius and $\tau$ is lower where the radiation density is lower. Thus, the observed temperature is set by the radiation energy density at the largest radius at which thermal equilibrium is kept. If, in contrast, the photons at the shock are out of thermal equilibrium, then the observed temperature can be much higher than $T_{\text{FB}}$ and is set by the number of photons that are generated over the diffusion time. This process is the same as in the case of a shock breakout from a stellar surface where no wind is present, and it is discussed in length in Katz et al. (2010) and NS10.

Assuming that photons and electrons in the unshocked wind have the same temperature, the observed temperature can be determined using the coupling coefficient $\eta \equiv 5 \times 10^{-3} E_{51}^{1/2} t_{\text{diff}}^{-1/2}$ defined in NS10, where all quantities are in cgs units. Here, $\eta$ is the ratio of the photon density needed to maintain thermal equilibrium to the photons produced per unit
of volume during the diffusion time, assuming that the electrons’ temperature is $T_{BB}$. This value of $\eta$ assumes that free–free emission dominates photon production. If bound–free opacity for $T_{BB}$ photons, $\kappa_{bf}$, is higher than free–free opacity, $\kappa_{ff}$, then $\eta$ is lower by the factor $\kappa_{bf}/\kappa_{ff}$ (see discussion in NS10). When the value of $\eta$ at the shock radius is $\eta_{n} < 1$, the radiation is in thermal equilibrium and the temperature is set by the electrons at the radius where $\eta = 1$.

At breakout, $aT_{BB}^{6/3} \approx 6(6/7)\eta_{w,bo}^{2/3} \eta_{bo}^{2}$, implying

$$T_{BB,bo} = \left(\frac{18c}{7\kappa_{ff}v_{bo,d}}\right)^{1/4} \approx 10^{8}v_{bo,d}^{-1/4} \text{K},$$

and yielding

$$\eta_{bo} \approx 5 \left(1 + \frac{\kappa_{bf}}{\kappa_{ff}}\right)^{-1} v_{bo,d}^{4/11},$$

where free–free processes are assumed to be the main source of photons (and absorption opacity). At any given time $\eta \propto \kappa^{11/8}$, implying that if $\eta_{bo} < 1$, then the observed temperature of the breakout pulse is $T_{BB,bo}^{6/11}$, yielding an observed breakout temperature

$$T_{bo,obs}(\eta_{bo} < 1) \approx 2 \times 10^{4} \left(1 + \frac{\kappa_{bf}}{\kappa_{ff}}\right)^{-6/11} v_{bo}^{2/11} \left(\frac{3000 \text{Km s}^{-1}}{v_{bo}}\right)^{2.2} t_{bo,d}^{-0.2} \text{K}.$$  

In contrast, when $\eta_{bo} \gg 1$, the observed breakout temperature is much higher than $T_{BB}$ and is given by $T_{BB,bo}^{6/11}$, yielding an observed breakout temperature

Figure 1 depicts the observed breakout temperature as a function of the breakout velocity for $t_{bo} = 20$ days. The solid line describes a wind composition such that photon production (and absorption) is dominated by free–free emission. In that case the critical breakout velocity, above which thermal equilibrium cannot be maintained, is $\approx 7000 \text{Km s}^{-1}$. It is evident that $T_{bo,obs}$ depends strongly on deviation from thermal equilibrium, rising by more than an order of magnitude when $v_{bo}$ is increased by a factor of two. Thus, for $v_{bo} = 0.3$ we find $T_{bo,obs} \approx 10^{8} \text{K}$, while if $v_{bo} = 1.5$, then $T_{bo,obs} \approx 5 \times 10^{9} \text{K}$. The effect of high metallicity is demonstrated by the dashed line in Figure 1, describing the observed breakout temperatures in the case that $\kappa_{bf}/\kappa_{ff} = 10$ (and $\eta_{bo}$ is smaller by that factor). Evidently, a higher metallicity implies a thermal equilibrium also at higher velocities. Note that $T_{bo,obs}$ at which thermal equilibrium breaks is almost independent of the exact opacity ($5 \times 10^{9} \text{K}$ in Figure 1). It also depends weakly on $t_{bo}$ and is always $<10^{9} \text{K}$. Since the velocity range that we consider here is exactly at the point where thermal equilibrium is marginal, we expect some of the breakouts in this range to be in equilibrium with $T_{bo,obs} \sim 10^{9} \sim 10^{10} \text{K}$, therefore bright in optical/UV, and some to be out of thermal equilibrium with $T_{bo,obs} > 10^{9} \text{K}$, bright in soft X-rays, but faint in optical.

3.3. Post-breakout Evolution

Following breakout, the collisionless shock that develops heats protons to high temperatures. The protons then transfer heat to electrons, which are cooled by radiation. The balance between the electron heating and cooling rates sets the electron temperature, during the post-shock cooling process (Katz et al.)
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The shock-heated electrons and protons are confined to a narrow layer behind the shock. Although they are the source of the observed luminosity, they do not necessarily set the typical photon temperature. This temperature is set during the diffusion of the photons toward the observer through the unshocked wind, as long as it is optically thick.\footnote{Fast cooling dictates that the shocked plasma is cooling, forming a dense cold layer behind the hot shock-heated plasma. This cold layer does not affect the observed soft component. The reason is that after the breakout the photons escape on a time that is shorter than the dynamical time, and therefore the energy density behind the shock is not dominated by radiation. Since the shock is radiative, the shocked plasma is compressed to keep the pressure balance, and once free–free cooling becomes dominant, its cooling rate increases with the compression and a runaway cooling decouples the plasma temperature from the rest of the system, preventing it from contributing a significant number of photons to the radiation field that cools the shock-heated plasma.}

The electrons of the unshocked wind are much colder than the \(\sim 60\) keV electrons immediately behind the shock. These unshocked electrons are heated by the diffusing radiation, and they produce photons that share energy with the diffusing radiation. During the early evolution after the breakout, the major part of the \(\sim 60\) keV electron layer cooling occurs as photons produced by the unshocked electrons diffuse back into the hot electron layer behind the shock and heat up by IC over this hot layer. These upscattered photons also serve during this stage as the main heating source of the unshocked electrons, through absorption of photons with similar temperature, and possibly also through Compton collisions with photons upscattered to \(m_e c^2/\gamma^2\), where Compton losses are substantial. The balance between the heating and the cooling of the unshocked electrons determines their temperature.

If unshocked electrons are able to produce a sufficient amount of photons during the available diffusion time, then the diffusing radiation is in thermal equilibrium. The condition for thermal equilibrium during the post-breakout evolution is similar to that during the breakout, namely, \(\eta < 1\). Thermal equilibrium also implies that each photon with \(T \sim T_{BB}\) is absorbed at least once during its diffusion through the unshocked electrons. This ensures that the unshocked electrons and the diffusing photons share a single temperature and the observed temperature is simply \(T_{BB}\) at the outermost layer that can maintain thermal equilibrium, namely, the layer where \(\eta = 1\). At breakout, the blackbody temperature of the diffusing radiation just ahead of the shock is \(T_{BB,s} = T_{BB,bo}\), and the thermal coupling coefficient of the electrons at this location is \(\eta_s = \eta_{bo}\). Following the breakout, \(T_{BB}\) and \(\eta_s\) (both measured at the shock radius) evolve as

\[
T_{BB,s}(t) \propto \left(\frac{L \tau}{r^2}\right)^{1/4} \propto t^{-3/4}, \tag{14}
\]

and

\[
\eta_s(t) \propto \frac{T_{BB,s}^{3.5}}{\rho^2} \propto T_{BB}^{3.5} r^{4} \propto \left\{\begin{array}{ll}
1 & n = 12 \\
0.6 & n = 7 \\
0.45 & n = 4
\end{array}\right. \tag{15}
\]

The evolution of the observed soft-component temperature \(T_{BB}\) where \(\eta = 1\) while thermal equilibrium is maintained is therefore

\[
T_{soft}(t, \eta_s < 1) \propto T_{BB,s}^{6/11} \eta_s^{6/11} \propto \left\{\begin{array}{ll}
0.2 & n = 12 \\
0.45 & n = 7 \\
1.1 & n = 4
\end{array}\right. \tag{16}
\]
Thus, while \( \eta_s < 1 \), the observed temperature decreases slowly at \( t < t_{SP} \) and drops linearly with \( t \) during the snowplow phase for as long as \( k_B T_{\text{soft}} \gtrsim 6000 \) K, beyond which recombination takes place and the temperature drop stops.

The main difference between the scenario discussed above (\( \eta_s < 1 \)) and the one in which the radiation ahead of the shock deviates from thermal equilibrium, i.e., \( \eta_s > 1 \), is that absorption of the photons that carry most of the luminosity by unshocked electrons stops playing a role. This has two important implications. First, the energy of a single photon can significantly grow by IC over the shocked hot electrons without being absorbed. The limit on the observed energy of such IC upscattered photons is set by three factors: Compton losses at \( m_ec^2/\tau^2 \), escape time from the system, and photoabsorption of extreme UV/soft X-rays by partially ionized pre-shock metals. Second, the coupling between the diffusing photons and unshocked electrons is done by Compton scattering only. This implies, in the parameter space that we consider, that the unshocked electrons and diffusing photons cannot maintain a single, well-defined temperature. As a result, the photons produce a non-trivial spectrum between the temperature of the hard shocked electrons and \( m_ec^2/\tau^2 \), and there may be no clear distinction between soft and hard components. Calculating the exact observed photon spectrum in that case is beyond the scope of this paper.

In general, our analysis of the soft component is not applicable anymore when the first of the following three takes place: (1) recombination becomes important, at which point \( T_{\text{soft}} \approx 6000 \) K; (2) the hard component carries most of the luminosity (which always precedes the inapplicability of the diffusion approximation due to low wind optical depth); and (3) \( \eta_s > 1 \).

### 3.3.2. Hard Component

The soft component is accompanied by harder photons, which can be generated by the hot electrons behind the shock in two ways, (1) free–free emission and (2) repeated IC upscattering of photons from the soft component. The total energy in IC upscattered photons is exponentially suppressed when \( \eta_s < 1 \), since only photons at the exponential tail of \( T_{\text{soft}} \) can be upscattered significantly before being absorbed by electrons with \( T_{\text{soft}} \). It may be suppressed farther at \( T \gg T_{\text{soft}} \) by photoabsorption of partially ionized metals ahead of the shock. We therefore consider here only hot electrons’ free–free emission. At breakout, the main cooling source of the hot electrons is IC of soft photons, yielding a ratio of free–free to IC emissivity (see Chevalier & Irwin 2012)

\[
\frac{\epsilon_{\text{ff,bo}}}{\epsilon_{\text{IC,bo}}} = \frac{a n_e (m_e c^2)^{3/2}}{4 k_B T_{\text{e,rad}}} \approx 10^{-2} v_{\text{bo,}\text{e},9}^{-2}, \tag{17}
\]

where \( \epsilon_{\text{rad}} \approx \rho v^2/2 \) at breakout. After the breakout, as long as \( \tau > 1 \),

\[
\frac{\epsilon_{\text{ff}}}{\epsilon_{\text{IC}}} (t) \propto \frac{1}{T^{1/2} v^3} \propto \begin{cases} t^{2.5} & t < t_{SP} \\ t > t_{SP} & \end{cases}. \tag{18}
\]

Therefore, IC remains the dominant cooling source for a long time after the breakout, suppressing the emission of hard photons. But even the hard photons that are emitted at \( T_e \) are not seen directly by the observer as long as the wind optical depth is significant. Compton losses during their diffusion toward the observer limit their energy to \( m_ec^2/\tau^2 \) and further reduce the luminosity of the hard component by a factor \( \tau^2 T_e/m_ec^2 \) (Chevalier & Irwin 2012). Therefore, the observed temperature of the hard photons is

\[
T_{\text{hard}}(t) = \min \left[ \frac{m_ec^2}{k_B \tau^2}, T_e(t) \right]. \tag{19}
\]

The luminosity of hard photons, counting only the contribution of free–free emission of the hot shocked electrons, is

\[
L_{\text{hard}}(t) \sim L(t) \frac{T_{\text{hard}}(t)}{T_{\text{e}}(t)} \min \left[ 1, \frac{\epsilon_{\text{ff}}}{\epsilon_{\text{IC}}}(t) \right]. \tag{20}
\]

Equations (13), (17), and (19) and \( t_{bo} = c/\nu_{\text{bo}} \) imply that when the breakout is at thermal equilibrium, this luminosity is always

\[
L_{\text{hard,bo}} \sim 10^{-4} L_{\text{bo}}. \tag{21}
\]

Equations (7) and (18) imply that at \( t < t_{SP} \), while \( L_{\text{hard}} \ll L_{\text{IC}} \), it first evolves rapidly and then slows down, within the range

\[
L_{\text{hard}}(t < t_{SP}) \propto t^{2.5} \to t. \tag{22}
\]

If the hard component is still subdominant during the transition to the snowplow phase, or if the breakout takes place during this phase, then its fraction out of the total luminosity rises quickly, and multiplying this fraction with \( L(t) \),

\[
L_{\text{hard}}(t > t_{SP}) \propto t^3 \to t^{0.5}, \tag{23}
\]

until it becomes comparable to the bolometric luminosity.

Finally, Equation (20) shows that right after breakout the hard photons from the hot electrons’ free–free emission carry only a tiny fraction (~10^{-4}) of the bolometric luminosity (which is seen in optical/UV). It implies that even though the contribution of repeated IC scattering of \( T_{\text{soft}} \) photons is strongly suppressed, it may still dominate the luminosity of the hard component. Therefore, while \( \eta_s \ll 1 \) implies \( L_{\text{hard,bo}} \ll L_{\text{bo}} \), the fraction may be higher than 10^{-4} right after the breakout.

### 3.4. Criterion for Fast Cooling

We assumed above that the shocked plasma is cooling fast (within a dynamical time). Here we find when this assumption holds. A sufficient requirement for fast cooling is that free–free emission cools the hot electrons over a dynamical time. In case the Compton \( \gamma \)-parameter drops below unity, it is also a necessary condition. The free–free cooling time at breakout is much shorter than the dynamical time for our considered range of \( \nu_{\text{bo}} \):

\[
\frac{\tau_{\text{cool,bo}}}{\tau_{\text{bo}}} \approx \frac{1}{25} v_{\text{bo,9}}^{-4} T_{e,60}^{-1/2}. \tag{24}
\]

Past breakout, it evolves as

\[
\frac{\tau_{\text{cool}}}{\tau} \propto v^{-2} T^{-1/2} \propto \begin{cases} t^{0.6} & n = 12 \\ t^{0.2} & n = 7 \\ t^{-0.5} & n = 4. \end{cases} \tag{25}
\]

Thus, if the shock is cooling fast at \( t_{SP} \), then it continues to cool fast also later. Requiring \( \tau_{\text{cool}}(t_{SP})/\tau_{SP} < 1 \) and taking \( \tau_{\text{cool}}/\tau \propto t^{0.6} (n = 12) \) all the way from breakthrough to \( t_{SP} \) shows that the shock enters the snowplow phase cooling fast as long as \( \nu_{\text{bo}} \gtrsim 6 v_{\text{bo,9}}^{12} M_{10}^{12} E_{51}^{-1/4} \) days. This calculation overestimates the minimal \( \nu_{\text{bo}} \) in Equation (25) by a factor of ~3 since \( n \) varies from \( n \lesssim 12 \) at breakout to \( n \approx 4 \) at \( t_{SP} \). Therefore, in the range of breakout velocities that we consider, the shock is always fast cooling for breakout times that are days or longer.
4. LIGHT CURVES OF THE DIFFERENT REGIMES

The light curve and spectral evolution depend on three timescales. The first two are $t_{\text{bo}}$ and $t_{SP}$, which affect the hydrodynamics. The third timescale, which we denote $t_{\text{hard}}$, marks the time at which the hard component becomes the dominant one, thus affecting the spectral evolution. For breakout times of a day or longer, the order of appearance of these three characteristic timescales depends primarily on the breakout time $t_{\text{bo}}$.

Since at breakout the hard component is always a small fraction of the total luminosity, $t_{\text{bo}} < t_{\text{hard}}$, the evolution depends on the location of $t_{SP}$ with respect to $t_{\text{bo}}$ and $t_{\text{hard}}$. One regime is $t_{SP} < t_{\text{bo}} < t_{\text{hard}}$. The criterion for this regime was found in Section 2. When $t_{\text{bo}} < t_{SP}$, there are two additional regimes, separated by the limiting case of $t_{SP} = t_{\text{hard}}$. In order to find the separating criterion, we use Equation (20). $t_{\text{hard}}$ is the earliest time that satisfies $L_{\text{hard}} \approx L$, namely, that both $\tau^2 \approx m_e c^2 / k_B T_e$ and $\kappa^2 \approx \kappa_{IC}$ are satisfied. Setting $t_{SP} = t_{\text{hard}}$, the condition $\tau^2 \approx m_e c^2 / k_B T_e$, which is the last to be satisfied for typical parameters, implies $\tau \approx 5E_{51}^{1/2} M_{10}^{-1/2}$, where we used $\tau(t_{SP}) \approx \sqrt{2E / M_{\text{ej}}}$ and Equation (13). Using Equation (5) and setting $\tau = 5E_{51}^{1/2} M_{10}^{-1/2}$, one finds that $t_{SP} = t_{\text{hard}}$ when $t_{bo,d} \approx 20M_{10}^{0.75} E_{51}^{-0.25}$. We thus obtain the following three regimes with respect to $t_{SP}$:

\begin{align*}
  t_{SP} < t_{\text{bo}} & \quad t_{bo,d} \gtrsim 80M_{10}^{0.75} E_{51}^{-0.25} \\
  t_{bo} < t_{SP} < t_{\text{hard}} & \quad 20M_{10}^{0.75} E_{51}^{-0.25} \lesssim t_{bo,d} \lesssim 80M_{10}^{0.75} E_{51}^{-0.25} \\
  t_{bo} < t_{\text{hard}} < t_{SP} & \quad 1 \lesssim t_{bo,d} < 20M_{10}^{0.75} E_{51}^{-0.25}.
\end{align*}

(26)

These regimes roughly match initial (breakout) $n$ values of 4, 7, and 12, respectively. Below we discuss the luminosity and spectral evolution in each of these regimes.

4.1. Late Breakout ($t_{SP} < t_{\text{bo}}$)

If the progenitor went through an extreme mass loss episode just prior to the explosion, the mass in the wind can be larger than the ejecta mass. In such a scenario the breakout may take place weeks to months after the explosion, near the time that, or even after, the reverse shock is exhausted. The entire explosion energy is released then around $t_{\text{bo}}$. If the ejecta is several solar masses with $\sim 10^{51}$ erg, the breakout velocity is 3000–5000 km s\(^{-1}\). At these velocities $\eta_s < 1$ throughout the evolution even for $\kappa_{hf} \lesssim \kappa_{hf}$, implying a very bright optical–UV flash, with a comparable rise time and full time and a typical temperature of $\sim (1-5) \times 10^9$ K. Following breakout, the luminosity and temperature of the soft component decline as $t^{-1.5}$ and $t^{-1.3}$, respectively. The temperature decline stops around 6000 K, where recombination starts playing an important role.

The hard component, which requires either a low ejecta mass $\sim M_{\odot}$ or a very high energy explosion $\sim 10^{52}$ erg, is expected in this regime only if the breakout velocity is high, which requires either a low ejecta mass $\sim M_{\odot}$ or a very high energy explosion $\sim 10^{52}$ erg. Figure 2 depicts the luminosity and temperature evolution of a breakout near the snowplow phase with $t_{\text{bo}} = 70$ days and $v_{\text{bo}} = 0.4$ matching $E_{51} = 1.65$ and $M_{10} = 10 M_{\odot}$), showing a very bright flash in the optical–UV, with a comparable rise time (not seen in the figure) and fall time. The temperature of the soft component (calculated for $\kappa_{hf} \lesssim \kappa_{hf}$ and the bolometric luminosity (solid lines) decline as $t^{-1.1}$ and $t^{-1.5}$, respectively. Temperature decline stops around 6000 K, where recombination starts playing an important role. The luminosity of hard photons that are generated by free–free emission of the hot shocked electron and their temperature (dashed red lines) are strongly suppressed around the breakout (see the text). $\gamma$eV photons are observed only more than a year after the explosion. The breakout time and velocity are chosen to fit the observation of SN 2006gy.

(A color version of this figure is available in the online journal.)

Figure 2. Breakout within the snowplow regime, with $t_{\text{bo}} = 70$ days and $v_{\text{bo}} = 0.4$ (matching $E_{51} = 1.65$ and $M_{10} = 10 M_{\odot}$), showing a very bright flash in the optical–UV, with a comparable rise time (not seen in the figure) and fall time. The temperature of the soft component (calculated for $\kappa_{hf} \lesssim \kappa_{hf}$ and the bolometric luminosity (solid lines) decline as $t^{-1.1}$ and $t^{-1.5}$, respectively.
Figure 3. Early breakout of $t_{bo} = 10$ days and $v_{bo} = 5000$ km s$^{-1}$, where $E_{SP}$ is calculated for $E_{51} = 1$ and $M_{10} = 1.25$. The bolometric luminosity and soft component temperature are marked with solid lines. The temporal decay of the bolometric luminosity is shifted from $t^{-0.3}$ ($n = 12$) to $t^{-0.6}$ ($n = 7$) at 0.1$t_{SP}$, to account for the flattening of the ejecta profile once $m$ approaches $M_{ej}$. Such a breakout may explain SNe like PTF 09uj (Ofek et al. 2010). $T_{soft}$ is plotted for two different absorption opacities. When $\kappa_{ff} \lesssim \kappa_{ff}$ (low metallicity), $T_{soft}$ remains in thermal equilibrium until the soft component becomes negligible. The luminosity and temperature of the hard photons are marked with dashed lines. The luminosity includes only the contribution from hot shocked electrons’ free–free emission, which is very faint at first, and becomes brighter and harder quickly to dominate the flux from day 200 on. Contribution from IC photons (not included in the hard component here) may be important at early time, especially in the case that $\kappa_{ff} \lesssim \kappa_{ff}$, where thermal equilibrium is marginal. This may result in significant X-ray emission also earlier than day 200. Such events make excellent candidates for X-ray searches ~100–300 days after the SN explosion and possibly even earlier. (A color version of this figure is available in the online journal.)

Figure 4. Intermediate regime breakout of $t_{bo} = 30$ days and $v_{bo} = 5000$ km s$^{-1}$. $E_{SP}$ is calculated for $E_{51} = 1$ and $M_{10} = 1$ and $T_{soft}$ calculated for $\kappa_{ff} \lesssim \kappa_{ff}$. The transition to the snowplow phase happens at day $\sim$230, while the hard component gains dominance only after this transition, reaching a lower luminosity and softer X-ray temperatures compared to those of the early regime. If the breakout is early enough, so that $T_{hard} \sim T_{SP}$, such events produce the brightest X-ray luminosity among all the cases. Notations are similar to Figures 2 and 3. (A color version of this figure is available in the online journal.)

Early breakout may explain SN PTF 09uj, as suggested by Ofek et al. (2010). The breakout temperature and luminosity of the case depicted in Figure 3 are consistent with the observations (Ofek et al. 2010).

4.3. Intermediate Breakout Time ($t_{bo} < t_{SP} < t_{hard}$): $20M_{10}^{0.75}E_{51}^{-0.25} \lesssim t_{bo,1} \lesssim 80M_{10}^{0.75}E_{51}^{-0.25}$

Figure 4 presents the luminosity and temperature of an intermediate case, where the breakout takes place before the transition to the snowplow phase, but the hard component may become dominant only after this transition. The properties of the breakout and the emission that follows after it are similar to those described in the early breakout case. The main difference from the early breakout case is that the hard component may not dominate the luminosity at $t_{SP}$, and by the time that it does, the luminosity may already be dropping fast. If the breakout is early enough and $t_{hard} \sim t_{SP}$, then the X-ray luminosity must be high. In fact, these events produce the brightest X-ray luminosity among all the cases discussed in this paper.

5. SUMMARY

We examine the emission from a shock breakout through a thick wind. We restrict our exploration to breakout timescales that are days or longer, which are easier to detect, and velocities in the range 3000–15,000 km s$^{-1}$, as expected for typical SNe. We further restrict our focus to cases where thermal equilibrium prevails during the evolution of the collisionless shock. For low metalicities this restriction reduces the shock velocities of interest to the lower part of the above range. We consider a standard wind profile $\rho_{w} \propto r^{-2}$, which is extended beyond the breakout radius. Following the breakout, the internal energy in these systems is dominated by the wind material that is heated by the forward shock (Chevalier & Irwin 2011; Balberg & Loeb...
2011), which makes a transition from an RMS to a collisionless shock at the breakout (Katz et al. 2011). Our main conclusions are listed below.

1. The breakout and following emission are very luminous. In the regime that we explore, the shocked plasma is always cooling fast (see the exact condition in Section 3.4), converting all the shock luminosity into radiation. Therefore, if the wind mass is comparable to, or larger than, the ejecta mass, then all the SN kinetic energy is radiated away. If the wind mass $M_w$ is smaller than $M_{ej}$, then the energy radiated during the interaction is $\sim E_{me}/M_w$. This result supports the suggestion by Chevalier & Irwin (2011) that at least some ultraluminous Type IIn SNe are such breakouts.

2. As long as photons that diffuse ahead of the shock maintain thermal equilibrium with the unshocked gas, the post-breakout emission is composed of two spectral components—soft and hard. The electrons behind the collisionless shock are heated to temperatures of $\sim \text{min}(60, 200 v_w^2)$ keV (Katz et al. 2011). The soft component is generated by the unshocked gas ahead of the shock, while the hard component is generated by the hot shocked electrons via free–free emission and IC of a small fraction of the soft photons.

3. The soft component may or may not be in thermal equilibrium during the breakout, even in the narrow velocity range of 3000–15,000 km s$^{-1}$. Thermal equilibrium is expected for all slower breakouts $\lesssim 7000$ km s$^{-1}$, and also for faster breakouts if there is a significant bound–free absorption (i.e., high metallicity). The breakout is then bright in optical/UV, with a temperature of $10^4–10^5$ K. In faster shocks that are out of thermal equilibrium, the breakout temperature can be as high as $5 \times 10^6$ K. In that case, the breakout X-rays can be much brighter than the optical/UV.

4. If thermal equilibrium is maintained ahead of the shock, the hard component is suppressed as long as IC cooling is more efficient than free–free cooling (Chevalier & Irwin 2012). In that case hard photons from free–free emission of hot electrons behind the forward shock carry at the breakout only $\sim 10^{-4}$ of the soft-component energy. The luminosity of hard photons from IC over the same electrons may be higher, but it is also a small fraction of the bolometric luminosity. The hard component luminosity rises quickly after breakout. It becomes dominant at $\sim 10–50 t_{bo}$.

5. The temperature of the hard component is $m_ec^2/\tau^2$, which at breakout is $0.1–1$ keV. It rises as $t^2$ after breakouts with $t_{bo} \lesssim 80$ days and as $t$ for longer breakouts.

The post-breakout evolution can be divided into three regimes according to the breakout time: (1) Late breakouts (typically $t_{bo} > 80$ days) are very bright in optical/UV and for typical parameters are very dim in X-rays at early time, $\sim 1$ keV X-rays peak around $20 t_{bo}$, carrying only $\sim 10^{-2}$ of the breakout energy. (2) Early breakout (typically 1 day $< t_{bo} < 20$ days) bolometric luminosity is lower and can be either in or out of thermal equilibrium for typical parameters. If it is out of thermal equilibrium, the emission can be bright in X-rays soon after the breakout (and possibly dim in optical/UV). If it is in thermal equilibrium, then X-rays are suppressed after the breakout, but their luminosity and hardness rise quickly to gain dominance, at most at $t \sim 20 t_{bo}$. (3) Intermediate breakouts (typically 20 days $< t_{bo} < 80$ days) are typically in thermal equilibrium and are therefore bright in optical/UV, while X-ray emission is suppressed at breakout. But $\sim 1–10$ keV X-rays become dominant at most at $t \sim 20 t_{bo}$, with a luminosity that may be almost as high as that of the breakout pulse.

In our calculations, we ignore the radius at which the thick wind ends. Once the wind density drops abruptly, the forward shock becomes inefficient and the luminosity fades quickly (Chevalier & Irwin 2011). Since the mass of the progenitor at birth is limited and the mass in the wind needed for an extended breakout is considerable, it is expected that the light curves we present will be terminated at some point. For example, in order for the light curve to remain bright until $t_{sp}$, the wind mass must be $\gtrsim M_{ej}$. Even when $M_w > M_{ej}$, the collected mass is $\propto t^{1/2}$ after $t_{sp}$, limiting the lifetime of the bright emission. In terms of the prospects of observing X-rays, this consideration makes earlier breakouts more attractive, since the mass that early breakout events collect by the time that the hard component dominates is lower.

We do not address the extinction of soft X-rays due to photoabsorption by partially ionized unshocked wind (Chevalier & Irwin 2012). This process may suppress softer-end photons of the hard component, as well as hard photons of the soft component, if the latter is far from thermal equilibrium. If photoabsorption is important, then the observed flux of soft X-rays is lower than the one that we predict.

Finally, our treatment ignores the contribution of emission from ejecta mass before it is shocked by the reverse shock. Once the optical depth of the wind drops, this emission becomes similar to the emission from a typical SN without a thick wind (e.g., typical IP, Ib, and Ic SNe). The emission from the forward shock that we consider here always outshines this emission (in terms of bolometric luminosity). But, at late times, when most of the forward shock emission is in hard photons, the contribution from inner layers may dominate the IR–optical emission.

To conclude, breakouts through thick winds produce very bright SNe with a potential for detection across the entire electromagnetic spectrum. The brightest phase is the breakout, which is typically dominated by optical–UV emission. In that case early X-ray emission is suppressed, but X-rays becomes harder and their luminosity rises quickly after the breakout, peaking at most after $\sim 10–50 t_{bo}$.

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REFERENCES

Arcavi, I., Gal-Yam, A., Yaron, O., et al. 2011, ApJ, 742, L18
Balberg, S., & Loeb, A. 2011, MNRAS, 414, 1715
Campana, S., Mangano, V., Blustin, A. J., et al. 2006, Nature, 442, 1008
Chevalier, R. A. 1982, ApJ, 258, 790
Chevalier, R. A., & Fransson, C. 2008, ApJ, 683, L135
Chevalier, R. A., & Irwin, C. M. 2011, ApJ, 729, L6
Chevalier, R. A., & Irwin, C. M. 2012, ApJ, 747, L17
Colgate, S. A. 1974, ApJ, 187, 333
Couch, S. M., Pooley, D., Wheeler, J. C., & Milosavljević, M. 2011, ApJ, 727, 104
Ensslin, L., & Burrows, A. 1992, ApJ, 393, 742
Falk, S. W. 1978, ApJ, 225, L133
Gezari, S., Dessart, L., Basa, S., et al. 2008, ApJ, 683, L131
Gezari, S., Rest, A., Huber, M. E., et al. 2010, ApJ, 720, L77
Imshennik, V. S., Nadezhin, D. K., & Utrobin, V. P. 1981, Ap&SS, 78, 105
Katz, B., Budnik, R., & Waxman, E. 2010, ApJ, 716, 781
Katz, B., Sapir, N., & Waxman, E. 2011, arXiv:1106.1898

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