X-Ray Light Curve and Spectra of Shock Breakout in a Wind

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Abstract

We investigate the properties of X-ray emission from shock breakout of a supernova in a stellar wind. We consider a simple model describing aspherical explosions, in which the shock front with an ellipsoidal shape propagates into the dense circumstellar matter. For this model, both X-ray light curves and spectra are simultaneously calculated using a Monte Carlo method. We show that the shock breakout occurs simultaneously in all directions in a steady and spherically symmetric wind. As a result, even for the aspherical explosion, the rise and decay timescales of the light curve do not significantly depend on the viewing angles. This fact suggests that the light curve of the shock breakout may be used as a probe of the wind mass-loss rate. We compare our results with the observed spectrum and light curve of X-ray outburst 080109/SN 2008D. The observation can be reproduced by an explosion with a shock velocity of 60% of the speed of light and circumstellar matter with a mass-loss rate of \(5 \times 10^{-4} M_\odot \text{yr}^{-1}\).

Key words: radiative transfer – shock waves – supernovae: general – supernovae: individual (SN 2008D)

1. Introduction

A core-collapse supernova (SN) emits a bright ultraviolet (UV)/X-ray flash, a so-called “shock breakout,” when photons generated from the shock escape upstream. Shock breakout has been studied for several decades, since Klein & Chevalier (1978) and Falk (1978). The timescale of the emission is determined by the light-crossing time of the radius of the source and the diffusion timescale of photons in the unshocked or shocked matter (e.g., Ensmann & Burrows 1992; Matzner & McKee 1999). Shock breakout is a powerful probe of the stellar radius and the structure of the outer layer of the star, since it should be associated with all core-collapse SNe, and the emission properties are highly sensitive to the behavior of the shock.

In 2008, the Swift/XRT accidentally detected X-ray outburst (XRO) 080109 (Soderberg et al. 2008), which was associated with a Type Ib SN, SN 2008D (Mazzali et al. 2008; Malesani et al. 2009; Modjaz et al. 2009; Tanaka et al. 2009a). The luminosity rapidly reached the maximum in the first \(\sim 100\) s and exponentially decayed until 600 s from the onset of the outburst. The peak luminosity and total radiated energy are \(6 \times 10^{43} \text{erg s}^{-1}\) and \(2 \times 10^{46} \text{erg} \), respectively. The Swift/XRT spectrum is well fitted by a power-law function, rather than a Planck function. Soderberg et al. (2008) also reported that a UV/optical emission was detected by the Swift/UVOT \(\sim 1\) day after XRO 080109, as well as a decreasing X-ray emission \(L = (1.0 \pm 0.3) \times 10^{39} \text{erg s}^{-1}\) in the energy range of 0.3–10 keV) by the Chandra X-ray Observatory \(\sim 10\) days after the Swift discovery. XRO 080109 and the subsequent fainter X-ray emission are believed to originate from shock breakout and interaction of the shock with circumstellar matter (CSM; Chevalier & Fransson 2008; Soderberg et al. 2008).

The timescale of XRO 080109 is closely related to the shock radius at the moment of breakout. When interpreting the rise time as the light-crossing time of the breakout radius, it must be \(\approx 10^{15}\) cm (Soderberg et al. 2008). Since it is larger than the typical radius of a Wolf–Rayet star, XRO 080109 is believed to originate from a dense CSM. The observed duration is consistent with the diffusion timescale of photons in the unshocked CSM, in which the shock breaks out at a radius of \((1.1–1.6) \times 10^{15}\) cm (Balberg & Loeb 2011). If the rise time is regarded as the shock expansion timescale, the shock radius is estimated to be \(\approx 6 \times 10^{14}\) cm (Svirski & Nakar 2014). Though the two estimated values are different by a factor of a few, they agree on the excess of the breakout radius compared to the typical radius of a Wolf–Rayet progenitor. In general, Wolf–Rayet stars blow winds with terminal velocities \(v_{\text{t}} \approx 1000\) km s\(^{-1}\) (Prinja et al. 1990; Hamann et al. 1995) at rates \(M\) in the range of \(10^{-5}\) to \(10^{-4}\) \(M_\odot \text{yr}^{-1}\) (Hamann et al. 1995; Nugis et al. 1998). Since a wind mass-loss event is known to play a significant role in the evolution of a massive star (Maeder & Meynet 1978; Meynet et al. 1994), studying shock breakout also enriches the understanding of massive star evolution shortly before the explosion. For this important reason, the properties (such as timescale and luminosity evolution) of emission from the shock breakout in a wind have been predicted by several theoretical studies (e.g., Balberg & Loeb 2011; Chevalier & Irwin 2011, 2012; Moriya et al. 2011; Svirski et al. 2012; Svirski & Nakar 2014).

The origin of the observed spectrum of XRO 080109 has been argued in several articles. The observed spectrum can also be fitted by a combination of two blackbody components, but the photospheric radii are far smaller than the typical radius of a Wolf–Rayet star (Li 2008). Soderberg et al. (2008) attributed the power-law spectral feature to electron (“bulk Comptonization”) scattering across a shock. In fact, Suzuki & Shigeyama (2010a) numerically examined how the photon energies increase due to this effect. Their results imply that the observed power-law X-ray spectrum requires a shock velocity higher than 0.3c, where c denotes the speed of light. The effect of bulk Comptonization has also been studied by Wang et al. (2007), in which mildly relativistic shock breakout in a dense CSM is applied for the low-luminosity GRB 060218/SN 2006j (Campana et al. 2006). Similar studies have been performed for shocks with lower velocities (\(< 10^3\) km s\(^{-1}\)) by Svirski et al. (2012) and Chevalier & Irwin (2012), applied to the luminous...
Type IIn SN 2006gy. The scattering process decreases photon energies in this particular SN.

In addition to the presence of the CSM and the bulk Comptonization, the asphericity of the shock front might also be important to determine the emission properties of shock breakout. Suzuki & Shigeyama (2010b) suggested that the shape of the light curve can reflect the degree of shock asphericity and the viewing angle. Couch et al. (2011) investigated the influence of shock asphericity on the light curve and spectrum by using results of their two-dimensional hydrodynamical simulations of a jet-driven SN. However, they do not take the influence of bulk Compton scattering into the calculation of the spectrum. Suzuki et al. (2016) recently performed 2D radiation hydrodynamic simulations for a blue supergiant exploding in a steady wind. Since bipolar explosions result in the shock appearing sequentially, the light curve would have a broader peak compared to the case of a spherical shock.

Despite a large number of studies having investigated the emission properties in detail, there have been no studies that reproduce both the observed X-ray spectrum and light curve by taking bulk Comptonization into account. In this paper, we aim to investigate the influence of shock asphericity and bulk Comptonization on the properties of emission from shock breakout in a wind. For this purpose, we perform radiative transfer calculations using a Monte Carlo method. In Section 2, we describe the settings for the shock and Monte Carlo calculation. In Section 3, we show the results and comparisons with the observed properties of XRO 080109/SN 2008D. In Section 4, we conclude this paper.

2. Methods

We calculate X-ray light curves and spectra of shock breakout emission in a dense CSM. In the following subsections, we describe our model for the propagation of the shock (Section 2.1) and the method to calculate radiative transfer (Section 2.2).

2.1. Model for Shock

To capture the properties of X-ray emission, we adopt a simple model of shock breakout in a wind as described below (Figure 1). Our calculation does not take into account the feedback from emission to the fluid motion.

The matter is radiation dominated (the adiabatic index $\gamma$ equals 4/3), and the radiation and matter are in thermal equilibrium below the photosphere. Here we focus on inverse Compton scattering in the shocked CSM of interest. For that, we think of SN ejecta as a piston and focus on modeling the forward shock propagating in the CSM. We ignore the presence of the shocked ejecta. Since the supposed ejecta density, $\approx 10^{19}$ cm$^{-3}$ at the moment of breakout (hereafter $t = t_0$, where $t$ denotes the time measured from the moment of explosion), is orders of magnitude higher than that of the shocked CSM ($\approx 10^{13}$ cm$^{-3}$), photons would be absorbed or scattered immediately at the contact surface. (Inward-traveling photons generated from the shock front can be blocked by a shell filled by ejecta with a uniform density and a mass of $\Delta M_{ej} \sim 2 \times 10^{-7} M_{\odot}$, the same order of magnitude as the total mass of the shocked CSM, $M_{ej} = 1 \times 10^{-7} M_{\odot}$, at $t = t_0$.)

Here $\Delta M_{ej}$ is estimated from the shell width $R_e - r_{\text{min}}$ of 1/$\kappa \rho_{ej} = 3 \times 10^4$ cm, where $\kappa = 0.2$ cm$^2$ g$^{-1}$ is the opacity for electron scattering, $R_e$ is the radius of the contact surface, and $\rho_{ej} = 1.8 \times 10^{-4}$ g cm$^{-3}$ is the mass density of ejecta. To see the influence of the structure behind the shock on the emission properties (shapes of the light curve and spectrum), we compare results of calculations with those using self-similar solutions of Chevalier (1982) with different density structures of the ejecta. As shown in Figure 13 (Appendix B), there is no significant difference in the shapes of the X-ray light curves between the models. We obtain spectra with similar shapes as long as the density of the ejecta has a steep slope ($n > 10$) as a function of radius (see Figure 14). Thus, three different regions (unshocked CSM, shocked CSM, and unshocked ejecta) are under consideration. Both the shocked CSM and ejecta move at constant velocities. The ejecta are assumed to have a uniform density and evolve in homologous expansion. The total mass is $10 M_{\odot}$.

We consider a shock having an elliptoidal shape. The shock radial velocity follows the formula

$$v(f, \theta) = \frac{1 - f}{[(1 - f)^{1/2} \cos^2 \theta + \sin^2 \theta]^2} \times v(f, 0),$$

(1)

where $f$ denotes the oblateness of the shock front and $\theta$ denotes the angle measured from the symmetric axis. If the kinetic energy of the ejecta is fixed, the shock velocity at $\theta = 0$ can be written as follows (Appendix A):

$$v(f, 0) = \sqrt{5} \times v_{f=0} \times (2f^2 - 4f + 3)^{-1/2},$$

(2)

where $v_{f=0}$ is the shock velocity in the spherically symmetric case. In this study, $v_{f=0} = 0.6c$. Figure 2 shows the angular dependence of $v$ for $f = 0$ (spherical), 0.1, 0.3, and 0.5.

The wind is supposed to be stationary, be spherically symmetric, and emanate from a carbon–oxygen layer. The electron number density $n_1$ of the unshocked CSM follows the equation

$$n_1 = \frac{A}{r^2},$$

(3)

where $r$ is the radius measured from the center of the progenitor and $A$ is a constant. The optical depth $\tau$ of the unshocked CSM must be equal to $c/v$ when the shock propagating at a speed $v$ breaks out. The characteristic timescale of the emission must

![Figure 1. Schematic view of an ellipsoidal shock propagating into a steady, spherically symmetric CSM. The degree of the shock asphericity is characterized by the oblateness $f$. The position of the shock front corresponds to the dashed curve at the moment of breakout.](image-url)

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**Figure 1.** Schematic view of an ellipsoidal shock propagating into a steady, spherically symmetric CSM. The degree of the shock asphericity is characterized by the oblateness $f$. The position of the shock front corresponds to the dashed curve at the moment of breakout.
strongly depend on $R_b$. We can determine the shock radius at the moment of breakout $R_b$ for $f = 0$ by the condition that the rise time of the observed emission is equal to the light-crossing time $R_b/c$. If we adopt $\Delta t_{\text{rise}} \approx 100$ s, then $R_b$ becomes $3 \times 10^{12}$ cm. Therefore, the constant $A$ can be uniquely determined by the following equation:

$$A = \frac{c}{\nu_f} \left[ \sigma_{\text{kl}} \int_{R_b}^{\infty} r^{-2} dr \right]^{-1},$$

with a free parameter $\nu_f=0$. Here $\sigma_{\text{kl}}$ is the Klein–Nishina cross section for a photon with an energy corresponding to the peak energy of a blackbody radiation. From the assumption of $\nu_f=0.6c$, $A = 7.5 \times 10^{36}$ cm$^{-1}$. When $f \neq 0$, $R_b$ has an angular dependence as written by

$$R_b = A\sigma_{\text{kl}} \times \frac{\nu(f, \theta)}{c}.$$

This equation indicates that the shock breaks out at the same moment $R_b/v$ in all directions (independent of $\theta$) if it has constant velocities. If the radial velocity of the unshocked CSM is 1000 km s$^{-1}$, the mass-loss rate $\dot{M}$ becomes $5 \times 10^{-4}\ M_\odot$ yr$^{-1}$. The rate is one order of magnitude higher than that for an ordinary Wolf–Rayet star, but still consistent with that of a luminous blue variable (Humphreys & Davidson 1994).

We can estimate the thickness $\Delta R$ of the shocked CSM assuming a uniform density $n_2$ there. The number density $n_2$ of the shocked CSM at the shock front satisfies the Rankine–Hugoniot relation

$$n_2 = n_1 \left( \frac{\gamma + 1}{\gamma - 1} \right).$$

A relation between the masses of the matter swept up by the shock and the shocked CSM is written by

$$4\pi R^2 \rho_2 \Delta R = 4\pi AR,$$

where $R = \nu(f, \theta) t$ is the shock radius. From Equations (6) and (7), $\Delta R$ equals $(\gamma - 1)/(\gamma + 1) \times R = R/7$. Therefore, the density $\rho_2 = 2.0 \times 10^{-11}$ g cm$^{-3}$ and the pressure $p_2 = 6.3 \times 10^9$ g cm$^{-1}$ s$^{-2}$ at the moment of shock breakout ($R = R_b$). Under the assumption of a radiation-dominated state and local thermo-dynamic equilibrium, the temperature for the shocked CSM is determined by $T_2 = (3\rho_2c^2)^{1/4} = 1.3 \times 10^6$ K and that for the ejecta by $T_{e2} = (3\rho_2c^2)^{1/4} = 6.9 \times 10^7$ K, where $\alpha$ is the radiation constant. The temperature of the unshocked CSM is $T_1 = 1.0 \times 10^4$ K, which is close to the typical effective temperature of a Wolf–Rayet star (Herald et al. 2000). The shocked CSM and the ejecta are assumed to have the same velocities $\nu(f, \theta)$. The assumption of $10\ M_\odot$ ejecta with high velocities of $\nu(f, \theta)$ itself is of course too energetic. Again, we note that photons do not enter a deep layer of the ejecta, so that only a very low mass ($\sim 2 \times 10^{-7}\ M_\odot$) of ejecta is required to have high velocities as $\nu(f, \theta)$. The kinetic energy of the ejecta in this region is $1 \times 10^{47}$ erg. For that reason, the supposed situation is not so bad in the region calculated in this work.

2.2. Monte Carlo Method

Using the settings of Section 2.1, we calculate radiative transfer of thermal photons by using a Monte Carlo method. The basic construction of the code is the same as we used in our previous study (Ohtani et al. 2013). Here we describe several assumptions made in this study.

Photons are isotropically generated at the shock front over a period of $\Delta t_{\text{ph}} = 0.5$ s. The period is determined so that the total radiation energy roughly equals the total emitted energy estimated from the Swift/XRT observation. A total of 1000 seed photons are generated every $5 \times 10^{-4}$ s, with an energy distribution following the Planck distribution in the rest frame of the fluid. If $f = 0$, the photospheric temperature is 0.11 keV (hereafter $k_B T_f = 0.6\ M_\odot$) at the moment of breakout, and the total energy $E_{\text{tot},i}$ radiated in the time interval $\Delta t_{\text{ph}}$ is $\sim 6 \times 10^{48}$ erg. After the shock breakout, some photons diffuse out of the shock front and reduce the pressure $p_2$ in the shocked CSM but do not significantly change the temperature, which is proportional to $p_2^{1/4}$. (From the thermal energy of the shocked CSM, $4\pi R^2 \Delta R\rho_2 T_2 = 1 \times 10^{48}$ erg, the change in the temperature $T_2$ is estimated to be $\sim 0.2\%$). Thus, we do not take into account this effect in the radiative transfer calculations.

We should note that the energy $E_{\text{tot},i}$ released by radiation is 1–2 orders of magnitude lower than the kinetic energy of the shocked matter, which must be comparable to the thermal radiation in an ordinary shock. We consider that this is because a major portion of photons generated in the ejecta remains trapped. To discuss that, a comparison of the dynamical timescale (hereafter $t_{\text{dyn}}$) of the shock, $R_b/0.6c \sim 170$ s, with the diffusion time of photons in the ejecta is needed. We have estimated the diffusion time $t_{\text{diff}}$ in a region between $r = r_{\text{min}}$ and $r = R$. If the flow expands linearly with time, the optical depth of the above region becomes unity when the shock reaches a radius (hereafter $R_r^\ast$) of $\sim 7 \times 10^{12}$ cm. Thus, the diffusion time $t_{\text{diff}}$ becomes $(7-3) \times 10^{12}$ cm$^2$/s $\sim 200$ s. Since $t_{\text{diff}}$ is longer than $t_{\text{dyn}}$, it seems that a major portion of photons are still trapped in the ejecta.

The generated photons are assumed to interact with matter via inverse Compton scattering and free–free absorption. In the shocked CSM, the effective optical thickness $\tau_r$ can be estimated by

$$\tau_r = \sqrt{\alpha^2(n_2^{\ast} + n_2 \sigma_{\text{kl}})} \Delta R,$$
where $\alpha^f$ is the absorption coefficient, due to free–free transition of electrons

$$\alpha^f = 3.7 \times 10^8 T^{-1/2} Z^2 n_0 \nu^{-3} \times (1 - e^{-\nu/k_BT}) \tilde{g}_{bb} \text{ cm}^{-1} \tag{9}$$

(Rybicki & Lightman 1979). $T$ is the temperature, $Z$ the atomic number, $n_0$ the number density of ion, $h$ the Planck constant, $\nu$ the frequency, and $\tilde{g}_{bb} \sim 1$ the Gaunt factor. If $f = 0$, $\tau = 4 \times 10^{-4} \ll 1$ for $\hbar \nu = 0.3 \text{ keV}$ and $Z = 8$ at the moment of shock breakout. Therefore, most photons are not absorbed by electrons.

The process of photon–electron coupling is discussed by Nakar & Sari (2010) and Katz et al. (2010) for shock breakout at a stellar surface and by Svirski et al. (2012) for that in a wind. Here we estimate the total number of thermal photons produced by bremsstrahlung emission. The total emissivity integrated over frequency is expressed by

$$\varepsilon^f = 1.4 \times 10^{-27} T^{1/2} n_e Z^2 \tilde{g}_{bb}, \tag{10}$$

where $n_e$ is the electron number density, $n_0$ the ion number density, $Z$ the electric charge of the ion, and $\tilde{g}_{bb} \sim 1$ the Gaunt factor (Rybicki & Lightman 1979). Assuming $n_e = n_{ej}$ ($n_{ej} = 6 \times 10^{19} \text{ cm}^{-3}$: electron number density in the ejecta), $T = T_{ej}$ and fully ionized oxygen gas, Equation (10) yields $\varepsilon^f = 3 \times 10^{-17} \text{ erg s}^{-1} \text{ cm}^{-3}$. Therefore, the energy $E^f$ radiated from the ejecta per unit time is roughly evaluated by

$$4\pi R_{ej}^2 (1/\kappa_{bb}) \varepsilon^f = 1 \times 10^{18} \text{ erg s}^{-1},$$

and the time required to release the energy of $E_{tot.}$ is $6 \times 10^3 \text{ s}$ (dividing $E^f$ by $3k_B T_{ej} = 0$, we can roughly estimate the number of photons as $2 \times 10^{39} \text{ s}^{-1}$). Since the required time is shorter than $\Delta t_{ph}$, we can consider that the generated photons are abundant enough so that the radiation and matter achieve thermal equilibrium.

Under the assumption of fully ionized gas, possible bound–free absorption is neglected. In order to show the validity of this assumption, we estimate the timescale for photoionization of oxygen in a process similar to that of Suzuki & Shigeyama (2010a). The bound–free cross section of OVI ions is written by

$$\sigma_{bf} = \left( \frac{64\pi n_{bf}}{3\sqrt{3}Z^2} \right) \alpha_\text{bf} \left( \frac{\chi}{\hbar \nu} \right)^3, \tag{11}$$

where $n = 1$ denotes the principal quantum number, $g_{bf} \sim 1$ the bound–free Gaunt factor, $\alpha_\text{bf}$ the fine-structure constant, $a_\text{b}$ the Bohr radius, and $\chi = Z^2 \alpha_\text{b}^2 m_e c^2/(2n^2) = 0.87 \text{ keV} (m_e$: electron mass) the ionization potential (Rybicki & Lightman 1979). For photons with energy $3k_B T_{ej} = 0$, Equation (11) yields $\sigma_{bf} = 2 \times 10^{-18} \text{ cm}^2 \gg \sigma_{af}$. Although this fact indicates that bound–free absorption is a dominant source of opacity, the interaction would not significantly affect the nonthermal component of the X-ray spectrum owing to a short timescale of photoionization. The timescale can be estimated by the total emitted energy and number of nonthermal X-ray photons. Using the luminosity of the thermal emission expressed by

$$L_{th} = 4\pi R_{ej}^2 \sigma_{SB} T_{ej}^4 = 2 \times 10^{46} \text{ erg s}^{-1} \tag{12}$$

($\sigma_{SB} = \alpha c/4$; the Stefan–Boltzmann constant) and the time interval $\Delta t_{ph}$, the total energy becomes

$$E_{th} = L_{th} \frac{e^{\Delta t_{ph}}}{v_{ej}} = 2 \times 10^{46} \text{ erg}. \tag{13}$$

Thus, the number density of photons with energies of a few keV is

$$n_{ph} = \varepsilon E_{th}/\text{few keV} \approx 10^{18} \epsilon \text{ cm}^{-3}, \tag{14}$$

where $\epsilon E_{th}$ (here $\epsilon$ is supposed to be $\sim 0.1$) is the total energy of nonthermal photons. Therefore, the timescale for bound–free absorption is

$$\tau_{bf} = \frac{1}{c \sigma_{bf} n_{ph}} \approx 10^{-9} \text{ s}. \tag{15}$$

Then we estimate the timescale of radiative recombination for fully ionized oxygen written by

$$\tau_{rad} = \frac{1}{\alpha_Z^f n^2}, \tag{16}$$

where

$$\alpha_Z^f = 5.197 \times 10^{-14} \beta^3/2$$

$$\times (0.4288 + 0.5 \ln \beta + 0.469 \beta^{-1/3}) \text{ cm}^3 \text{s}^{-1}, \tag{17}$$

with $\beta = \chi/(k_B T)$ (Seaton 1959). Substituting $T = T_{ph}$ into the equations above, $\tau_{rad}$ becomes 0.6 s, which is far longer than $\tau_{bf}$.

We can expect from a comparison of the estimated $\tau_{rad}$ and $\tau_{bf}$ that most elements become fully ionized immediately, so that photons with energies of a few keV would not be influenced by bound–free absorption in the shocked CSM. In the unshocked CSM, we obtain $\tau_{bf} \approx 10^{-7} \text{ s}$ and $\tau_{rad} = 0.4 \text{ s}$ by replacing $\Delta t_{ph}$ by the diffusion timescale of photons of $\sim 100 \text{ s}$ (calculated in Section 3.1) in Equation (14), $T$ by $T_{ej}$, and $n^2$ by $n_1$ at the shock front in Equations (16) and (17). The small $\tau_{bf}/\tau_{rad}$ ratio allows neglecting the influence of bound–free transition.

Once a photon reaches the surface with an optical depth $\tau$ of $10^{-2}$, it is supposed to escape from the CSM. The calculation stops when the shock front reaches the surface of $\tau = 10^{-2}$.

3. Results

First of all, we calculate the X-ray light curve and spectrum for a spherically symmetric SN ($f = 0$) and compare them with the observation of XRO 080109. Then we show the dependence on the oblateness $f$ of the shock and the viewing angle $\Theta$.

3.1. Spherically Symmetric Shock

Figure 3 shows the light curve in the energy range from 0.3 to 10 keV covered by the Swift/XRT. The luminosity rapidly increases for the first 40 s and exponentially decreases for the subsequent several hundred seconds, as $L \propto \exp[(t_{ph}-t_{peak})/\tau]$, where $t_{peak} = 40 \text{ s}$ and $\tau = 200 \text{ s}$. The time interval between the onset and the peak (hereafter $\Delta t_{onset}$) depends primarily on the light-crossing time $\Delta t_{sc}$ of the size of the emerging shock and secondarily on the diffusion timescale of photons in the shocked CSM. We should note that there is a weak but not negligible
The time is measured from the moment when the first photon passes a large spherical surface concentric with the ejecta.

The emission in the energy range of the XRT rise and the exponential decay resembles that of the observed XRO 080109. Though the XRT spectrum can be reproduced by the emission generated from a spherically symmetric shock with a velocity of $0.6c$ and a wind with a mass-loss rate of $5 \times 10^{-4} M_{\odot} \text{yr}^{-1}$. Here the total radiation energy $E_{\text{tot,f}}$ is $2 \times 10^{46}$ erg, about 3 times higher than that before electron scattering ($E_{\text{tot,i}}$). The total kinetic energy $E_{\text{sh}}$ of the shocked CSM is $4\pi R_0^2 \rho_2 \Delta R v^2 = 2 \times 10^{47}$ erg at the moment of shock.
breakout. From the relatively small ratio of $E_{\text{tot},f}-E_{\text{tot},i}$ to $E_{\text{sh}}$, we can expect that radiation feedback would not induce a significant change of electron temperature.

3.2. Aspherical Shock

We investigate the influence of the asphericity of the shock on the light curve and the spectrum. Previous calculations (Suzuki & Shigeyama 2010b; Couch et al. 2011; Matzner et al. 2013; Salbi et al. 2014; Suzuki et al. 2016) investigated aspherical shock breakout in the vicinities of the stellar surfaces without thick CSM. Caused by the significant time lags between the shocks breaking off their tops and sides, the calculated light curves show broader peaks compared to that for a spherical shock. The situation of our calculation is fairly different from that, as the shock breakout occurs simultaneously in all directions owing to the assumptions of the thick, steady, and spherically symmetric wind and a constant shock velocity $v(f, \theta)$. We should note that, in reality, nonradial motions of ejecta play important roles along the stellar surface. If the asphericity (or “obliquity”) in the ejecta motion is limited to a thin outer layer of the star and the effect of radiation is neglected, nonradial flows are believed to suppress the shock (Matzner et al. 2013). Thus, the assumptions in our model need a somewhat energetic explosion process, such as a jet-like explosion or prolonged activity of the central engine.

Figure 6 shows the light curves in the energy range of the Swift/XRT when the shock has an oblateness of 0.1, 0.3, and 0.5. The time $t_{\text{obs}}$ is measured from the moment when the first photon passes a large spherical surface concentric with the ejecta. In each panel, the flux is averaged over the angular ranges of $\theta = [0^\circ, 10^\circ], [40^\circ, 50^\circ], \text{ and } [80^\circ, 90^\circ]$, respectively. The light curves show that the timescales (the rise time and the duration) of the luminosity evolution have similar values regardless of $f$ and $\theta$. As with the spherically symmetric case, the decay time of a few hundred seconds is uniquely determined by the density distribution of the CSM. The reason for the similarity in the rise time is rather complicated. It depends on the value of the shock velocity, which varies with the inclination angle. When the radiation intensity is concentrated in a small angle ($<\theta_{\text{rad}}$), the rise time $t_{\text{rise}}$ becomes significantly shorter than the light-crossing time $\Delta t_{\text{lc}}$ of the size of the emerging shock, whereas when the intensity is broadly distributed ($v = 0.5c$ in Figure 7), $\Delta t_{\text{rise}}$ roughly equals $\Delta t_{\text{lc}}$. For example, we can estimate the rise time observed with a viewing angle of $\Theta = 0$ for a shock with an oblateness of $f = 0.5$. The high shock velocity along the line of sight of $\sim 0.8c$ implies that most photons reaching an observer travel close to the symmetry axis ($\theta < 30^\circ$) owing to beaming effects (see Figure 7). Thus, the rise time is approximated as $[R_{b,\theta=0} - R_{b,\theta=30^\circ}]c \sim 40$ s, where $R_{b,\theta=0} = 4 \times 10^{12}$ cm and $R_{b,\theta=30^\circ} = 3 \times 10^{12}$ cm. On the other hand, the rise time observed with a viewing angle $\Theta = 90^\circ$ can be approximated by $\Delta t_{\text{lc}} = 40$ s because of the low shock velocity along the line of sight of $\sim 0.5c$.

The peak luminosity $L_{\text{peak}}$ decreases with $\Theta$, as well as the velocity of the shock propagating along the line of sight. The relation between the shock velocity $v$ and $L_{\text{peak}}$ in the energy range of 0.3–10 keV is displayed in Figure 8. The bin widths of $v$ correspond to the ranges of the viewing angle $\Theta$. Within the range of $0.4c \lesssim v \lesssim 0.8c$, the values of $L_{\text{peak}}$ tend to increase with increasing $v$ owing to the bulk Comptonization. For a
lower shock velocity, such a positive correlation would become weaker because photons cannot gain much energy from electrons via electron scattering. Figure 9 shows the light curves for energy lower than 0.3 keV. Again, all of the models have similar results in the timescales and the overall evolution of the luminosity owing to the simultaneous shock breakouts in all directions. Thus, the timescales responsible for the shape of the light curve become independent of the oblateness of the shock and the viewing angle. Consequently, from such a light curve it would be possible to know whether the morphology of the CSM is spherical and the velocity of the shock propagating along the line of sight. On the other hand, it would be difficult to constrain the degree of asphericities of the ejecta and shock front.

Figure 10 shows the time-integrated spectra for \( f = 0.1, 0.3, \) and 0.5. When the shock front has a finite oblateness, the power-law gradient of the high-energy (1–7 keV) component is shallower along the on-axis compared to the spherically symmetric case, while it is much steeper along the off-axis. This is because the shock has a radial velocity higher than 0.6c in the vicinity of the axis and lower at off-axis. Consequently, the influence of bulk Comptonization becomes weaker as \( \Theta \) becomes larger.

We compare the calculation for the ellipsoidal shock wave with XRO 080109. Figure 11 shows the light curve for \( f = 0.5 \) and \( \Theta = 45^\circ \) as our best model. The overall shape of the light curve is roughly consistent with observations. In fact, an off-axis line of sight has also been suggested for XRO 080109 from late-phase observations of nebular emission lines of SN 2008D (Tanaka et al. 2009b), and our results are consistent with this interpretation. As shown in Figure 10, the high-energy spectral gradient is within a 1\( \sigma \) error range of the observation if the shock velocity along the line of sight is higher than \( \sim 0.5c \).

4. Conclusions

We investigate the properties of X-ray emission from shock breakout in a dense CSM. For this purpose, we calculate transfer of X-ray photons interacting with matter through Compton scattering and free–free absorption by using a Monte Carlo method. We also study relations between the asphericity of the shape of the ellipsoidal shock front and the observational features of the emission.

The rise time of the light curve \( \Delta t_{\text{rise}} \) is mainly determined by the light-crossing time \( t_{\text{lc}} \) of the breakout radius and also slightly affected by the diffusion time \( t_{\text{diff}} \) of photons in the shocked CSM. The major factor determining the duration of the light curve is the light-crossing time of the radius at which photons last scatter off electrons. Even for an aspherical explosion, the properties of the light curve, such as the duration, the rise time, and the shape of the declining part, do not dramatically depend on the viewing angle as long as a steady and spherically symmetric wind is considered. The result suggests that the characteristics of the light curve are a good probe of the CSM density or mass-loss rate.

We show that both the observed light curve and spectrum of XRO 080109/SN 2008D can be reproduced by mildly relativistic shock breakout in a dense spherical CSM. For a shock with a velocity of 0.6c and CSM with a mass-loss rate of \( 5 \times 10^{-4} M_{\odot} \) yr\(^{-1} \), the rise time, the duration, and the shape of the calculated X-ray light curve can be consistent with the observation. The power-law spectral gradient of the observed emission is also reproduced if the shock propagates toward the observer at a speed greater than \( \sim 0.5c \).
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Appendix A

Relation of Shock Velocity to Oblateness

The shape of the shock front should be changed under a fixed value of the explosion energy $E_{\text{kin}}$. In the following few equations, we approximate the velocity of ejecta by that of the shock front. Here we define a typical time $t_0$ for each model, at which the mass density $\rho_0$ of ejecta has the same value as the moment of shock breakout for $f = 0$. Parameter $t_0$ equals $(1 - f)^{-2/3} t_s$, where $t_s$ denotes the moment of shock breakout along the symmetric axis. Then $E_{\text{kin}}$ and the total ejecta mass $M_{\text{ej}}$ are roughly expressed by the following equations:

$$E_{\text{kin}} \approx \iint \int \rho v^2 f \sin \theta d\theta d\phi,$$

$$= \frac{2\pi}{3} \rho_0 v_0^2 R_0^3 (2f^2 - 4f + 3)(1 - f)^2,$$

where $\phi$ is the azimuth angle, $R$ is the shock radius, and $\rho_0$, $v_0$ and $R_0$ are the values at $\theta = 0$. Accordingly,

$$E_{\text{kin}} \propto v_0^2 \times (2f^2 - 4f + 3).$$

As a result, with a fixed $v_0 = 0.6c$, $v_0$ can be written by Equation (2).

Appendix B

Collision of the Spherically Symmetric Ejecta with Power-law Density and the Circumstellar Matter

In order to see whether the assumption of uniform density significantly influences the results of the calculation, we make a calculation for spherically symmetric ejecta with a power-law density ($\propto r^{-n}$). The issue of the absolute luminosity is beyond the scope of this work. We describe the hydrodynamics by using the Chevalier self-similar solution. We examine the emission for $n = 10$, 12, and 7 (ordinarily used to describe SNe Ia).

B.1. Settings

To calculate the X-ray light curve and spectra, we describe the spherically symmetric distribution of the shocked matter and freely expanding ejecta as follows.

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[Figure 10. Time-integrated spectra of X-ray emission originated from an axisymmetric shock. The dashed curve represents a blackbody spectrum at a temperature of $kTB$. The gradients of the straight lines are $-2.0$ and $-2.6$, which correspond to the ±1σ values of those of the observed Swift/XRT spectrum.]

[Figure 11. Same as Figure 3, but for an ellipsoidal shock with an oblateness of $f = 0.5$.]
The structure of the shocked CSM is determined by Parker (1963), in which the radius of the forward shock $R_1$ increases with time $t$ as $R_1 \propto t^{1/\lambda_{\text{r}}}$ ($\lambda_{\text{r}}$: constant). In the stationary CSM, if $\lambda_{\text{r}} = 3/2$, the total energy of the shock becomes constant. Here we assume the same density profile of the unshocked CSM and the shock radius $R_1$ at the moment of breakout as those ($n_1$ and $R_{\text{bc},f=0}$) in Section 2.1.

The structure of the shocked ejecta is derived by Chevalier (1982), in which the density of the unshocked ejecta is assumed by

$$\rho_{\text{ej}} = r^{-3} \left( \frac{r}{L_{\text{g}}} \right)^{-n}$$

(g, $n$: constant) and the radius $R_2 \propto t^{1/\lambda_{\text{r}}}$ of the reverse shock, where $\lambda_{\text{r}} = (n - 2)/(n - 3)$. Noting that the pressure and velocity of the shocked matter should be continuous at the contact surface, $g$ becomes $4.6 \times 10^9$ for $n = 7, 7.6 \times 10^9$ for $n = 10$, and $9.0 \times 10^9$ for $n = 12$. Figure 12 displays the fluid profile.

Using the hydrodynamical profile above, the X-ray light curves and time-integrated spectra for the fixed shock velocity $0.6c$ at the moment of shock breakout are calculated. The settings for the Monte Carlo calculation are the same as those in Section 2.2, with the exception of the following. If the velocity of the forward shock is fixed at $0.6c$ at the moment of shock breakout, $t_b$ becomes 133 s for $n = 7$, 146 s for $n = 10$, and 150 s for $n = 12$. The temperature of the matter is $64 \text{ eV}$ (hereafter $k_{\text{B}}T_{2,\text{ch}}$) at the shock front. From the electron number density $n_{\text{ej}}$ of $\approx 10^{14} \text{ cm}^{-3}$ in the ejecta and temperature $T$ of $T_{2,\text{ch}}$, the total emissivity of free–free emission $\varepsilon_{\text{ff}}^r$ is estimated to be $\approx 10^7 \text{ erg s}^{-1} \text{ cm}^{-3}$ and the energy $E_{\text{ff}}^r$ radiated from the ejecta per unit time is $\approx 10^{33} \text{ erg s}^{-1}$. Therefore, it takes $\approx 10^3 - 10^4 \text{ s}$ to release the energy of $E_{\text{tot}}$. This fact would not lead to the conclusion that it is impossible to generate photons by free–free emission, but a more careful study would be needed to discuss the structure of the shock in detail.

B.2. Light Curves and Spectra

We investigate how the shapes of the light curve and X-ray spectrum differ from those calculated in Section 3. For example, if the density of the shocked matter follows a uniform distribution, the rise time $\Delta t_{\text{rise}}$ strongly depends on the light-crossing time $\Delta t_{\text{lc}}$ of the breakout radius and weakly on diffusion time $t_{\text{diff}}$. Here we investigate the dependence of the emission properties on the structure of the shock in the spherically symmetric case.

Figure 13 shows the resultant light curves in the energy range of 0.3–10 keV. The graphs are compared with the model with uniform density distribution (black dotted line; the luminosity is scaled by a factor of 0.05). If $n = 10$ and 12, the overall shapes of the light curves are quite similar to that for the model with uniform density. In order to know on what timescales the rise time $\Delta t_{\text{rise}}$ and duration depend, we compare the diffusion time $t_{\text{diff}} = (R_{p=1} - r_{\text{min}})/\nu$ in the shocked matter and unshocked CSM with the light-crossing time $\Delta t_{\text{lc}}$ of the size of the emerging shock. Here $R_{p=1} \approx 7 \times 10^{12} \text{ cm}$. From the radius $r_{\text{min}}$ (which satisfies $\int_{r_{\text{min}}}^{R_p} \kappa \rho dr = 1$), $2.58 \times 10^{12} \text{ cm}$ for $n = 10$ and $2.61 \times 10^{12} \text{ cm}$ for $n = 12$, $t_{\text{diff}}$ is estimated to be $\approx (R_{p=1} - r_{\text{min}})/0.5 \approx 300 \text{ s}$, which is longer than the light-crossing $\Delta t_{\text{lc}}$ estimated in Section 3.1. For this reason, $\Delta t_{\text{lc}}$ can be said to be the primary factor in determining $\Delta t_{\text{rise}}$ rather than $t_{\text{diff}}$.

Though a model with $n = 7$ is fainter than those with $n = 10$ and $n = 12$ owing to the low temperature of the shocked matter, the timescale of the brightening and the duration are not significantly different from those with $n = 10$ and 12. (The radius $r_{\text{min}}$ equals $2.45 \times 10^{12} \text{ cm}$, so that $t_{\text{diff}} > \Delta t_{\text{lc}}$.)
Figure 14 shows the time-integrated spectrum. Here we examine how the spectral gradient of the high-energy (1–7 keV) component changes with the motion of the matter. In comparison with the model with the uniform density and velocity distribution (black short-dashed line), the spectral gradient is steeper if the structure behind the shock is considered. This is because the velocity of the shocked matter is lower than the shock wave. For Figure 14, we can say that the spectral gradient of the high-energy tail is determined by the effect of bulk Comptonization, so it can be a source of information on the shock velocity.

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References
Balberg, S., & Loeb, A. 2011, MNRAS, 414, 1715
Campana, S., Mangano, V., Blustin, A. J., et al. 2006, Natur, 442, 1008
Chevalier, R. A. 1982, ApJ, 258, 799
Chevalier, R. A., & Fransson, C. 2008, ApJL, 683, L135
Chevalier, R. A., & Irwin, C. M. 2011, ApJL, 729, L6
Chevalier, R. A., & Irwin, C. M. 2012, ApJL, 747, L17
Couch, S. M., Pooley, D., Wheeler, J. C., & Milosavljević, M. 2011, ApJ, 727, 104
Ensman, L., & Burrows, A. 1992, ApJ, 393, 742
Falk, S. W. 1978, ApJL, 225, L133
Hamann, W.-R., Koesterke, L., & Wessolowski, U. 1995, A&A, 299, 151
Herald, J. E., Schulte-Ladbeck, R. E., Eenens, P. R. J., & Morris, P. 2000, ApJS, 126, 469
Humphreys, R. M., & Davidson, K. 1994, PASP, 106, 1025
Katz, B., Budnik, R., & Waxman, E. 2010, ApJ, 716, 781
Klein, R. I., & Chevalier, R. A. 1978, ApJL, 223, L109
Li, L.-X. 2008, MNRAS, 388, 603
Maeder, A., & Meynet, G. 1987, A&A, 182, 243
Malesani, D., Fynbo, J. P. U., Jhorth, J., et al. 2009, ApJL, 692, L84
Matzner, C. D., Levin, Y., & Ro, S. 2013, ApJ, 779, 60
Matzner, C. D., & McKee, C. F. 1999, ApJ, 510, 379
Mazzali, P. A., Valenti, S., Della Valle, M., et al. 2008, Sci, 321, 1185
Meynet, G., Maeder, A., Schaller, G., Schaefer, D., & Charbonnel, C. 1994, A&AS, 103, 97
Modjaz, M., Li, W., Butler, N., et al. 2009, ApJ, 702, 226
Moriya, T., Tominaga, N., Blinnikov, S. I., Baklanov, P. V., & Sorokina, E. I. 2011, MNRAS, 415, 199
Nakar, E., & Sari, R. 2010, ApJ, 725, 904
Nugis, T., Crowther, P. A., & Willis, A. J. 1998, A&A, 333, 956
Ohtani, Y., Suzuki, A., & Shigeyama, T. 2013, ApJ, 777, 113
Parker, E. N. 1963, Interplanetary Dynamical Processes (New York: Interscience Publishers)
Prinja, R. K., Barlow, M. J., & Howarth, I. D. 1990, ApJ, 361, 607
Rybicki, G. B., & Lightman, A. P. 1979, Radiative Processes in Astrophysics (New York: Wiley)
Salbi, P., Matzner, C. D., Ro, S., & Levin, Y. 2014, ApJ, 790, 71
Seaton, M. J. 1959, MNRAS, 119, 81
Soderberg, A. M., Berger, E., Page, K. L., et al. 2008, Natur, 453, 469
Suzuki, A., Maeda, K., & Shigeyama, T. 2016, ApJ, 825, 92
Suzuki, A., & Shigeyama, T. 2010a, ApJ, 719, 881
Suzuki, A., & Shigeyama, T. 2010b, ApJL, 717, L154
Svirski, G., & Nakar, E. 2014, ApJL, 788, L14
Svirski, G., Nakar, E., & Sari, R. 2012, ApJ, 759, 108
Tanaka, M., Tominaga, N., Nomoto, K., et al. 2009a, ApJ, 692, 1131
Tanaka, M., Yamanaka, M., Maeda, K., et al. 2009b, ApJ, 700, 1680
Wang, X.-Y., Li, Z., Waxman, E., & Meszaros, P. 2007, ApJ, 664, 1026