SHOCK BREAKOUT DRIVEN BY THE REMNANT OF A NEUTRON STAR BINARY MERGER: AN X-RAY PRECURSOR OF MERGERNOVA EMISSION

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ABSTRACT

A supra-massive neutron star (NS) spinning extremely rapidly could survive from a merger of an NS-NS binary. The spin-down of this remnant NS that is highly magnetized could power the isotropic merger ejecta to produce a bright mergernova emission in the ultraviolet/optical bands. Before the mergernova, the early interaction between the NS wind and the ejecta could drive a forward shock propagating outward into the ejecta. As a result, a remarkable amount of heat can be accumulated behind the shock front and the final escape of this heat could produce a shock breakout emission. We describe the dynamics and thermal emission of this shock with a semi-analytical model. It is found that a few hours after the merger, by leading the mergernova emission as a precursor, sharp and luminous breakout emission appears mainly in soft X-rays, with a luminosity of $\sim 10^{45}$ erg s$^{-1}$. Therefore, the detection of such an X-ray precursor could provide evidence for identifying NS-powered mergernovae and distinguishing them from radioactivity-powered novae (i.e., kilonovae or macronovae). The discovery of NS-powered mergernovae could finally help to confirm the gravitational wave signals due to the mergers and the existence of supra-massive NSs.

Key words: gamma-ray burst; general – stars: neutron – supernovae: general

1. INTRODUCTION

A merger of binary neutron stars (NSs) is one of the most promising targets for directly detecting gravitational waves (GWs). Such a detection, pointing to a few hundred mega-parsecs, is expected to come true in the very near future via the second generation of ground-based GW detectors (Abadie et al. 2010; Niitsumae et al. 2013). To meet this expectation, the electromagnetic counterparts of the GW signals must necessarily be observed, helping to confirm the position, time, redshift, and astrophysical properties of the GW sources.

During an NS-NS merger event, the rapidly rotating compact system could eject a highly collimated relativistic jet and a quasi-isotropic sub-relativistic outflow, where various electromagnetic emission could be produced. Specifically, internal and external dissipations of the jet energy could result in a short-duration gamma-ray burst (GRB) and its multi-wavelength afterglow emission, which could be the most attractive electromagnetic counterparts, considering their high brightness (see Nakar 2007; Berger 2014 for reviews). However, unfortunately, the detection probability of an associated GRB is significantly suppressed by the small opening angle of the jet (Fong et al. 2014), although a late and weak orphan afterglow could still be observed off-axis (van Eerten & MacFadyen 2012; Metzger & Berger 2012). Alternatively, more and more attention has been paid to the emission originating from the isotropic ejecta, e.g., thermal emission due to the diffusion of internal energy of the ejecta (called mergernova; Li & Paczyński 1998; Kulkarni 2005; Rosswog 2005; Metzger et al. 2010; Yu et al. 2013) and non-thermal emission due to the shock interaction between the ejecta and the environment medium (Nakar & Piran 2011; Metzger & Berger 2012; Gao et al. 2013; Piran et al. 2013).

The isotropic merger ejecta is probably highly neutron-rich, which makes it possible to effectively synthesize nuclei much heavier than $^{56}$Ni via rapid neutron capture processes (r-processes; Rosswog et al. 1999, 2014; Goriely et al. 2011; Roberts et al. 2011; Korobkin et al. 2012; Bauswein et al. 2013a; Hotokezaka et al. 2013, Hotokezaka & Piran 2015; Wanajo et al. 2014; Goriely 2015; Just et al. 2015; Martin et al. 2015). The radioactive decays of these newly synthesized elements can heat the ejecta to produce a detectable thermal emission. However, due to limitations imposed by the low mass of the ejecta (no more than a few percent of solar mass) and the element synthesis efficiency, the luminosity of a radioactive-powered mergernova is expected to not be much higher than $\sim 10^{41}$ erg s$^{-1}$. Thus, these phenomena are widely known as kilonovae or macronovae (Kulkarni 2005; Metzger et al. 2010; Barnes & Kasen 2013; Tanaka & Hotokezaka 2013; Grossman et al. 2014; Kasen et al. 2015a).

This luminosity limit can be breached if a more powerful energy source is provided by the central post-merger compact object. By invoking a remnant supra-massive NS that spins extremely rapidly and is highly magnetized, Yu et al. (2013) and Metzger & Piro (2014) investigated the characteristics of mergernova emission using an energy injection from the spin-down of the NS. It was found that the peak luminosity of such an NS-powered mergernova, which of course depends on the lifetime of the NS and the collimation of the NS wind, could sometimes be comparable to, and even exceed, the luminosity of ordinary supernovae, but with a much shorter duration on the order of a few days. Exceedingly, some unusual rapidly evolving and luminous transients were discovered by some recent observations, such as the Pan-STARRS1 Medium Deep Survey (Drout et al. 2014). The characteristics of these transients are basically consistent with the predictions for NS-powered mergernovae (Yu et al. 2015), although multi-wavelength cross-identifications are still needed.

Recently, some multi-wavelength studies have been carried out on observational candidates of mergernovae. For example, as implied from the shallow-decay afterglows of GRB 130603B, the widely discussed infrared excess after this GRB has been argued to probably be powered by a remnant NS (Fan et al. 2013) rather than by radioactivity, as usually...
considered (Berger et al. 2013; Tanvir et al. 2013). More directly, by considering the high-energy emission from an NS wind (Yu et al. 2010; Zhang 2013), which can partly leak from a preceding merger ejecta at late time, an NS-powered mergernova is expected to possibly be accompanied by a late X-ray re-brightening. Very recently, Gao et al. (2015) found that the late optical and X-ray bumps after GRB 080503 provided a perfect example of such multi-wavelength features. In this paper, we shall reveal another possible X-ray signature prior to an NS-powered mergernova. This signature is caused by the breakout of the shock arising from the early interaction between the NS wind and the merger ejecta. This X-ray precursor emission could play an essential role in future mergernova identifications and in GW detections.

2. THE MODEL

2.1. Remnant NS and Merger Ejecta

A merger would inevitably happen in an NS-NS binary once the gravity between the NSs could no longer be supported by angular momentum because of GW radiation release. After the merger, in some situations, a supra-massive NS rather than a black hole could be left with an extremely rapid differential rotation. This is supported indirectly by many emission features of short GRBs and afterglows, including extended gamma-ray emission, X-ray flares, and plateaus (Dai et al. 2006; Fan & Xu 2006; Rowlinson et al. 2010, 2013; Giacomazzo & Perna 2013; Zhang 2013). The presence of such a supra-massive NS is permitted by both theoretical simulations (Hotokezaka et al. 2011; Bauswein et al. 2013b) and observational constraints (e.g., the present lower limit of the maximum mass of Galactic NSs is precisely set by PSR J0348+0432 to 2.01 ± 0.04M⊙; Antoniadis et al. 2013). On one hand, during the first few seconds after the birth of a remnant NS, a great amount of neutrinos can be emitted from the very hot NS. On the other hand, differential rotation of the NS can generate multipolar magnetic fields through some dynamo mechanisms (Duncan & Thompson 1992; Price & Rosswog 2006; Cheng & Yu 2014), making the NS a magnetar. Consequently, during the very early stage, the neutrino emission and ultra-strong magnetic fields together could drive a continuous baryon outflow (Dessart et al. 2009; Siegel et al. 2014; Siegel & Ciolfi 2015a), which provides an important contribution to form an isotropic merger ejecta. A few seconds later, the NS eventually enters into a uniform rotation stage and meanwhile a stable dipole structure of magnetic fields can form. From then on, the NS starts to lose its rotational energy via a Poynting flux-dominated wind. The spin-down luminosity carried by the wind can be estimated with the magnetic dipole radiation formula as \( L_{\text{d}} = L_{\text{ad}}(1 + t/\tau_{\text{md}})^{-2} \), where \( L_{\text{ad}} = 10^{45}R_s^6B_s^2\rho^{-1}_{\text{l-3}} \text{ erg s}^{-1} \), \( \tau_{\text{md}} = 2 \times 10^5R_s^6B_s^{-2} \rho_{\text{l-3}}^{-1} \text{ s} \), and the zero point of time \( t \) is set at the beginning of the magnetic dipole radiation, which is somewhat later than the NS formation by several seconds. Here \( R_s \), \( B_s \), and \( \rho_{\text{l-3}} \) are the radius, dipolar magnetic field strength, and initial spin periods of the NS, respectively, and the convention \( Q_s = Q/10^4 \) is adopted in cgs units.

During a merger event, although the majority of matter finally falls into the central remnant NS, there is still a small fraction of matter ejected outward, e.g., the baryon wind blown from the NS mentioned above. Besides that component, a quasi-isotropic merger ejecta can also be contributed by a wind from a short-lived disk surrounding the NS, by an outflow from the colliding interface between the two progenitor NSs, and by a tidal tail due to the gravitational and hydrodynamical interactions. The latter two components are usually called dynamical components. These ejecta components differ from each other in mass, electron fraction, and entropy. It is difficult and nearly impossible to precisely describe the specific constitutions and distributions of a merger ejecta, which depend on the relative sizes of the two progenitor NSs, the equations of state of the NSs’ matter, and magnetic field structures. In any case, according to the numerical simulations in the literature, on one hand, the mass of the dynamical components can range from \( \sim 10^{-4}M_\odot \) to a few times \( \sim 0.01M_\odot \) (Oechslin et al. 2007; Bauswein et al. 2013a; Hotokezaka et al. 2013; Rosswog 2013). On the other hand, in the presence of a remnant supra-massive NS, the mass of a neutrino-driven wind is found to be at least higher than \( 3.5 \times 10^{-3}M_\odot \) (Perego et al. 2014), and the mass-loss rate due to ultra-high magnetic fields is about \( 10^{-3} - 10^{-2}M_\odot \text{ s}^{-1} \) during the first 1–10 s (Siegel et al. 2014). Therefore, by combining all of these contributions (see Rosswog 2015 for a brief review), we can take \( M_d = 0.01M_\odot \) as a reference value for the total mass of an ejecta. Furthermore, we can adopt power-law density and velocity distributions of this mass as follows (Nagakura et al. 2014):

\[
\rho_{\text{ej}}(r, t) = \frac{(\delta - 3)M_d}{4\pi r_{\text{max}}^2} \left( \frac{r_{\text{min}}}{r_{\text{max}}} \right)^{3-\delta} - 1 \left( \frac{r}{r_{\text{max}}} \right)^{-\delta},
\]

and

\[
v_{\text{ej}}(r, t) = v_{\text{max}} \frac{r}{r_{\text{max}}(t)}, \text{ for } r \leq r_{\text{max}}(t),
\]

where \( r \) is the radius to the central NS, \( v_{\text{max}} \) is the maximum velocity of the head of the ejecta, which is probably on the order of \( \sim 0.1c \), and the slope \( \delta \) ranges from 3 to 4 according to the numerical simulation of Hotokezaka et al. (2013). We fix \( \delta = 3.5 \) as in Nagakura et al. (2014). In fact, the variation of \( \delta \) within a wide range cannot significantly affect the primary results of this paper except for an extremely high value (e.g., \( \delta > 5 \)). The maximum radius of ejecta can be calculated by \( r_{\text{max}}(t) \approx r_{\text{max},i} + v_{\text{max}}(t) \), where \( r_{\text{max},i} \approx v_{\text{max}} \Delta t, \) with \( \Delta t \) being the time on which the dipolar magnetic field is stabilized. Correspondingly, the minimum radius reads \( r_{\text{min}}(t) = r_{\text{min},i} + v_{\text{min}}(t) \) and its initial value can be determined by an escape radius as \( r_{\text{min}}(0) = (2GM_\odot/v_{\text{max}}^2)^{1/3} \), where \( G \) is the gravitational constant and \( M_\odot \) is the mass of the remnant NS.

The huge energy released from a remnant NS (i.e., millisecond magnetar) could eventually drive an ultra-relativistic wind mixing Poynting flux and leptonic plasma, which catches up with the preceding merger ejecta very quickly. On one hand, if this wind always remains Poynting-dominant, even until it collides with the ejecta, the material at the bottom of ejecta could be heated by absorbing low-frequency electromagnetic waves from the Poynting component. On the other hand, more likely, some internal dissipations (e.g., the ICMPRT processes; Zhang & Yan 2011) could take place in the NS wind to produce non-thermal high-energy emission. Subsequently, a termination shock could be formed at the interface between the wind and ejecta, if the wind
magnetization has become sufficiently low there (Mao et al. 2010). As a result, the bottom of the ejecta can be heated by absorbing high-energy photons from the emitting wind region and/or by transmitting heat from the neighbor hot termination shock region. Additionally, even if we arbitrarily assume an extreme situation in which all of the wind energy is completely reflected from the interface, the material at the bottom of the ejecta would also be heated due to adiabatic compression by the high pressure of the wind. Therefore, in any case, the energy carried by the wind can always be mostly injected into the bottom of ejecta and heat the material there.

2.2. Shock Heating and Emission

When the bottom of a merger ejecta is heated by an injected NS wind, a pressure balance can naturally be built between the wind and the ejecta bottom. On one hand, the pressure balance can gradually extend to larger radii through thermal diffusion, which, however, happens very slowly for an extremely optical thick ejecta. On the other hand, the high pressure of the bottom material can lead itself to expand adiabatically and to reach a high speed. This speed could result in the formation of a forward shock propagating outward into the ejecta.

By denoting the radius and speed of a shock front by \( r_{sh} \) and \( v_{sh} \), respectively, the increase rate of the mass of shocked ejecta can be calculated by

\[
\frac{dM}{dt} = 4\pi r_{sh}^2 \rho_{sh} \left[ v_{sh} - v_{j}(r_{sh}, t) \right],
\]

where \( v_{j}(r_{sh}, t) \) and \( \rho_{sh}(r_{sh}, t) \) are the velocity and density of the upstream material just in front of the shock. Obviously, we also have \( dM_{sh} = v_{sh} dt \). As the shock propagates, its bulk kinetic energy, which has been previously gained from adiabatic acceleration, can again be partly converted into the internal energy of the newly shocked material. The rate of this shock heating effect can be written as

\[
H_{sh} = \frac{1}{2} \left[ v_{sh} - v_{j}(r_{sh}, t) \right]^2 \frac{dM}{dt}.
\]

Then the total internal energy accumulated by the shock, \( U_{sh} \), can be derived from

\[
\frac{dU_{sh}}{dt} = H_{sh} - P_{sh} \frac{d(\epsilon V)}{dt} - L_{sh}.
\]

Here \( P_{sh} = U_{sh}/(3\epsilon V) \) is an average pressure, \( V \sim (4/3)\pi r_{sh}^3 \) is the volume of the whole shocked region experiencing an adiabatic expansion, and the fraction \( \epsilon \) is introduced by considering that the shock-accumulated heat is mostly deposited at a small volume immediately behind the shock front. The product of this average pressure and the corresponding volume represents the cooling effect due to adiabatic expansion. \( L_{sh} \) is the luminosity of shock thermal emission, which is caused by the diffusion and the escape of the shock heat.

Approximately following a steady diffusion equation

\[
L = \frac{4\pi \rho^2}{3\kappa \rho} \left( \frac{\partial u}{\partial r} \right) c, \quad \kappa = \text{opacity}, \quad \rho \text{ and } u \text{ are densities of mass and internal energy, respectively, we roughly estimate the luminosity of shock thermal emission by}
\]

\[
L_{sh} \approx \frac{\varepsilon_{sh}^2 U_{sh} c}{c r_{sh}^3 + \left( \frac{\varepsilon_{max}}{\varepsilon_{sh}} - r_{sh}^3 \right) \left[ 1 - e^{-\left( \varepsilon_{sh} + \tau_{un} \right)} \right]},
\]

where \( \tau_{sh} = \kappa M/4\pi r_{sh}^2 \) and \( \tau_{un} = \int_{r_{sh}}^{r_{un}} \kappa \rho dr \) are the optical depths of the shocked and unshocked ejecta, respectively. Here, the former optical depth, which only influences the decrease of the shock emission after breakout, is calculated by considering that most of the shocked material is concentrated within a thin shell behind the shock front (Kasen et al. 2015b). The value of the parameter \( \epsilon \) can be fixed by equating the shock luminosity during the breakout to the simultaneous heating rate, because after that moment freshly injected shock heat can escape from the ejecta almost freely. The opacity of the merger ejecta is predicted to be on the order of magnitude of \( \sim (10–100) \text{ cm}^2 \text{ g}^{-1} \), which results from the bound–bound, bound–free, and free–free transitions of lanthanides synthesized in the ejecta (Kasen et al. 2013). This value is much higher than the typical value of \( \kappa = 0.2 \text{ cm}^2 \text{ g}^{-1} \) for normal supernova ejecta. In this paper we take \( \kappa = 10 \text{ cm}^2 \text{ g}^{-1} \). Fairly speaking, some reducing effects on the opacity could exist, e.g., (1) the lanthanide synthesis in the wind components of ejecta can be blocked by neutrino irradiation from the remnant NS by enhancing electron fractions (Metzger & Fernández 2014) and (2) the lanthanides in the dynamical components could be ionized by the X-ray emission from the NS wind (Metzger & Piro 2014).

2.3. Shock Dynamics

The temporal evolution of shock thermal emission, as presented in Equation (6), is obviously dependent on the dynamical evolution of the shock. Due to the slow thermal transmission in the optical thick ejecta, a significant pressure/temperature gradient must exist in the ejecta during shock propagation, which could lead the shock to be continuously accelerated. Therefore, the dynamical evolution of such a radiation-mediated shock is completely different from the internal and external shocks in GRB situations. For an optical thin GRB ejecta, a pressure balance can be built throughout the whole shocked region simultaneously with the shock propagation. In that case, the shock velocity can simply be derived from shock jump conditions (Dai 2004; Yu & Dai 2007; Mao et al. 2010). On the contrary, in the present optical thick case, a detailed dynamical calculation of the shock in principle requires an elaborate description of the energy and mass distributions of the ejecta (see Kasen et al. 2015b for a one-dimensional hydrodynamical simulation for a similar process), which is beyond the scope of an analytical model. Nevertheless, a simplified and effective dynamical equation can still be obtained according to the energy conservation of the system.

The total energy of the shocked region can be written as \( E = \frac{1}{2} M_{sh} v_{sh}^2 + U \), where \( U \) is the total internal energy of the shocked region. For a radiation-mediated shock, the value of \( U \) should be much higher than \( U_{sh} \). In principle, the concept of a "shocked region" that is used here can generally include the NS wind regions, because the total mass of wind leptons is drastically smaller than that of shocked ejecta, and more importantly, the energy released from the NS is continuously distributed in both the wind and ejecta through thermal diffusion, which makes them behave like a whole. By ignoring the relatively weak energy supply from radioactivity, the
variation of the total energy can be written as

$$dE = (\xi L_{sd} - L_c)dt + \frac{1}{2} \nu_{\xi}^2 (r_{sh}, t) dM,$$  \hspace{1cm} (7)

where \( \xi \) represents the energy injection efficiency from the NS wind, which could be much smaller than one if the NS wind is highly anisotropic, and \( L_c \) is the total luminosity of the thermal emission of merger ejecta. As a general expression, the specific form of the energy injection is not taken into account. Substituting the expression of \( E \) into Equation (7), we can find the dynamical equation of the forward shock to be

$$\frac{dv_{sh}}{dt} = \frac{1}{M_{vsh}} \left[ (\xi L_{sd} - L_c) - \frac{1}{2} (v_{sh}^2 - \nu_{\xi}^2) \frac{dM}{dt} - \frac{d} {dt} \right].$$  \hspace{1cm} (8)

In order to clarify the expressions of \( L_c \) and \( dU/dt \), note that \( \dot{U} = U - U_{sh} \), which represents the internal energy excluding the shock-accumulated part. The evolution of this internal energy component can be given by

$$\frac{dU}{dt} = \xi L_{sd} - \bar{P} \frac{dV}{dt} - L_{mn},$$  \hspace{1cm} (9)

where adiabatic cooling is also calculated with an average pressure as \( \bar{P} = \dot{U}/3V \) and

$$L_{mn} \approx \frac{U_c}{t_{max}} \left[ 1 - e^{-(\tau_a + \tau_m)} \right].$$  \hspace{1cm} (10)

The above expression is different from Equation (6) because the majority of the internal energy of the shocked region is deposited in the innermost part of the region, which is much deeper than the shock front. Then we have \( L_c = L_{sh} + L_{mn} \). The emission component represented by \( L_{mn} \) actually is the merger nova emission discussed in Yu et al. (2013).

By substituting Equations (4), (5), and (9) into (8), we can obtain another form of the dynamical equation as

$$\frac{dv_{sh}}{dt} = \frac{1}{M_{vsh}} \left( \frac{U}{3V} \right) \frac{dV}{dt},$$  \hspace{1cm} (11)

which is just the expression adopted in Kasen et al. (2015b) to calculate the shock breakout for super-luminous supernovae. This equation can be easily understood in the framework of adiabatic acceleration of a “fireball.” As discussed above, what is accelerating the ejecta material is actually the local internal pressures at different radii, which are much lower than the pressure of the NS wind, due to the significant delay of pressure transmission. The work performed by these varying pressures can be effectively estimated with an average internal pressure of \( (U/3V) \) with respect to a volume variation of \( dV = 4\pi r_{sh}^2 v_{sh} dt \). With the above dynamical equation, the energy conservation and assignments of the system can be described reliably. By considering that the internal energy at the time \( t \) is on the order of magnitude \( U \sim L_{sd} t \), Equation (11) can naturally determine a kinetic energy that is also on the same order of magnitude \( \frac{1}{2} M_{vsh}^2 v_{sh} \sim L_{sd} t \).

3. RESULTS AND ANALYSES

A supra-massive NS surviving from a merger event is believed to initially spin with a Keplerian limit period of about 1 ms. This corresponds to a rotational energy of several times 10^{52} \text{ erg}, with a high stellar mass, which could be much higher in view of the rapid differential rotation. Most of this energy is probably consumed very quickly in order to generate and amplify magnetic fields, to drive a short GRB, and maybe also to radiate GWs. The duration of this violent stage should not be much longer than the duration of the short GRB. So we would take \( \Delta t = 2 \text{ s} \) which is the boundary dividing long and short GRBs. When a steady magnetic dipole radiation begins, the spin period could have been reduced to a few milliseconds, corresponding to an energy of \( \sim 10^{51-52} \text{ erg} \). This energy supply could be further discounted for the merger ejecta by the parameter \( \xi \), if a remarkable fraction is collimated within a small cone to power an extended gamma-ray emission and X-ray afterglow plateau after the GRB.

For typical values of spin-down luminosity and timescale as \( \xi L_{sd} = 10^{47} \text{ erg s}^{-1} \) and \( t_{rad} = 10^4 \text{ s} \), we present a representative numerical result in Figure 1. As shown in the top panel, the shock initially moves slowly, with a velocity much smaller than the maximum velocity of ejecta, and experiences a gradual acceleration process. When the shock velocity exceeds the maximum ejecta velocity by about a few tens of percentage, shock breakout happens, as indicated by the sharp peak in the middle panel. It takes a remarkably long time (\( \sim 10^3 \text{ s} \)) for the shock to breakout from the ejecta. This period is much longer than the shock breakout time gained in some previous works (\( \sim 1-10 \text{ s} \); Gao et al. 2013; Siegel & Ciolfi 2015a, 2015b;
Wang et al. (2015) because there the shock velocity was significantly overestimated by using shock jump conditions with an assumed global pressure balance. However, as discussed above, such a global pressure balance actually cannot be built very quickly for a radiation-mediated shock.

The internal energy produced by the shock only occupies a very small fraction of the total internal energy behind the shock. The shock jump conditions could be satisfied only after a very long time acceleration or after the ejecta was close to optically thin, before which the shock has already crossed the whole ejecta. The middle panel of Figure 1 also shows that the shock breakout emission is very luminous, which is comparable to that of the following bright mergernova emission peaking at a few days. Nevertheless, since the shock breakout and mergernova are emitted at very different radii, the corresponding emission temperature (∼10^4 K) of the former can be much higher than the latter (∼10^3 K), as presented in the bottom panel. Here an effective temperature is defined by \( T_e = (L_e / 4 \pi r_{\text{max}}^2 \sigma)^{1/4} \) with \( \sigma \) being the Stefan–Boltzmann constant. In more detail, we plot three chromatic light curves in Figure 2, assuming a blackbody spectrum7

\[
\nu L_\nu = \frac{8 \pi^2 r_{\text{max}}^2}{h^3 c^2} \frac{(\nu h)^4}{\exp(h\nu/kT_e) - 1},
\]

with temperatures given in the bottom panel of Figure 1, where \( h \) is the Planck constant and \( k \) is the Boltzmann constant. As shown, while the mergernova emission falls into the ultraviolet band, for the adopted model parameters, the shock breakout is mainly concentrated within soft X-rays. Finally, in order to verify some of the energy arguments mentioned above, we plot the temporal evolutions of different energy components in Figure 3. It is shown that, although the energy released from the NS is initially injected into the ejecta in the form of internal energy, the majority of this internal energy is finally converted into the kinetic energy of the ejecta. As a result, during the whole optically thick period, the internal and kinetic energies basically remain comparable with each other. The internal energy produced by the shock is obviously less than the injected one, but at the shock breakout time, the instantaneous release of this small amount of energy can still temporarily overshadow the emission component due to the thermal diffusion.

For a straightforward understanding of the characteristics of the shock breakout emission, here we provide some analytical analyses. First, in physics, shock breakout happens when the dynamical time begins to be longer than the diffusion timescale on which photons diffuse from the shock front to the outermost surface of merger ejecta. Hence we can in principle solve the shock breakout time \( t_{\text{bo}} \) from the equation

\[
t_{\text{bo}} = t_d = \left( \frac{r_{\text{max}} - r_{\text{sh,bo}}}{\lambda} \right)^2 \frac{\lambda}{c} = \left( \frac{r_{\text{max}} - r_{\text{sh,bo}}}{\kappa \rho_{\text{ej}}} \right)^2 \frac{\kappa \rho_{\text{ej}}}{c},
\]

where \( \lambda = 1/\kappa \rho_{\text{ej}} \) is the average path of photons. Approximately, we get

\[
r_{\text{sh,bo}} \approx r_{\text{max}} \left( 1 - \frac{r_{\text{max}}}{r_{\text{h}}} \right),
\]

where

\[
r_{\text{h}} \approx (v_{\text{max}} \kappa M_{\text{ej}}/c)^{1/2} = 4.5 \times 10^{15} \text{ cm} \quad (v_{\text{max}}/0.1 c)^{1/2} (\kappa/10 \text{ cm}^2 \text{ g}^{-1})^{1/2} (M_{\text{ej}}/0.01 M_{\odot})^{1/2}.
\]

This means that when the shock breakout happens, the shock radius has been very close to the maximum radius of ejecta. By invoking \( U \sim \xi L_{\text{sd}}, \quad P = U/(3V) \), and an acceleration rate \( a \sim 4 \pi r_{\text{sh,bo}}^2 P/M, \) we can estimate the breakout radius by

\[
r_{\text{sh,bo}} \sim \frac{1}{2} \left( \frac{\xi L_{\text{sd}}}{2 M_{\text{ej}}} \right)^{1/2} \left( \frac{3}{2} \right)^{1/2}, \quad \text{where } M \approx M_{\text{ej}} \text{ is adopted}
\]
because \( r_{sh,bo} \approx r_{max} \). Then, from the equation \( r_{sh,bo} \approx r_{max} = v_{max} t_{bo} \), we can simply derive the shock breakout time to
\[
T_{bo} \sim \frac{2M_{ej}v_{max}^2}{\xi L_{sd}}
\]
\[
= 3600 s \left( \frac{v_{max}}{0.1 c} \right)^2 \left( \frac{M_{ej}}{0.01 M_\odot} \right) \left( \frac{\xi L_{sd}}{10^{47} \text{ erg s}^{-1}} \right)^{-1}.
\]
Furthermore, we get \( v_{sh,bo} \sim 2v_{max} \) and \( r_{sh,bo} \sim 2M_{ej}v_{max}^3/\xi L_{sd} = 1.1 \times 10^{13} \text{ cm} \), which is indeed much smaller than \( r_g \). Then the luminosity and the temperature of shock breakout can be estimated to
\[
L_{sh,bo} \sim H_{sh} \sim 2\pi r_{sh,bo}^2 (v_{sh,bo} - v_{max})^3 \rho_{ej}
\]
\[
\sim 8 \times 10^{45} \text{ erg s}^{-1} \left( \frac{\xi L_{sd}}{10^{47} \text{ erg s}^{-1}} \right),
\]
and
\[
T_{sh,bo} = \left( \frac{L_{sh,bo}}{4\pi r_{max}^2 \sigma} \right)^{1/4} \sim 6 \times 10^5 \text{ K}
\]
\[
\times \left( \frac{v_{max}}{0.1 c} \right)^{-3/2} \left( \frac{M_{ej}}{0.01 M_\odot} \right)^{-1/2} \left( \frac{\xi L_{sd}}{10^{47} \text{ erg s}^{-1}} \right)^{3/4}.
\]

The above analytical expressions qualitatively exhibit the physical mechanisms and the parameter dependencies of the shock breakout emission, although the numbers given here are somewhat higher than the numerical ones because a linear acceleration is assumed. It is indicated that the dependence of the shock breakout emission on the density profile of the merger ejecta (i.e., the parameter \( \delta \)) is very weak. In more detail, in Figure 4 we present the variations of the shock breakout luminosity \( L_{sh,bo} \) and the peak photon energy \( \epsilon_p = 4kT_{sh,bo} \) in the \( \xi L_{sd} - t_{rad} \) parameter space. It is indicated that, for a significantly bright shock breakout emission, a large amount of energy is required to be released from a NS within a sufficiently short time.

4. CONCLUSION AND DISCUSSIONS

Mergernovae, particularly the novae powered by a remnant supra-massive NS, are some of the most competitive electromagnetic counterparts of GW signals during NS-NS mergers. The discovery of NS-powered mergernovae could substantially modify and expand our conventional understandings of supernova-like transient phenomena. Therefore, understanding how to identify mergernovae is an essential question for future searches and observations. Undoubtedly, a multi-wavelength method is necessary and would be helpful for answering this question. In this paper, we discover that a shock breakout can be driven by the early interaction between a merger ejecta and a succeeding NS wind. Such a breakout appears a few hours after the merger by leading the mergernova emission as a precursor. The breakout emission would mainly be in soft X-rays with a luminosity of \( \sim (10^{44} - 10^{46}) \text{ erg s}^{-1} \), corresponding to an X-ray flux of a few \( \times (10^{-11} - 10^{-9}) \text{ erg s}^{-1} \text{ cm}^{-2} \) for a distance of \( \sim 200 \text{ Mpc} \), which can be above the sensitivity of many current and future telescopes, e.g., the Swift X-ray telescope (Burrows et al. 2005), Einstein Probe (Yuan et al. 2015), etc. More optimistically, some X-ray shock breakout emission could have appeared in some X-ray afterglows of short GRBs, probably exhibited as early X-ray flares. It will be interesting to discover such candidates.

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