Inverse Compton Scattered Merger-nova: Late X-Ray Counterpart of Gravitational-wave Signals from NS–NS/BH Mergers

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Abstract

The recent observations of GW170817 and its electromagnetic (EM) counterparts show that double neutron star mergers could lead to rich and bright EM emissions. Recent numerical simulations suggest that neutron star and neutron star/black hole (NS–NS/BH) mergers would leave behind a central remnant surrounded by a mildly isotropic ejecta. The central remnant could launch a collimated jet and when the jet propagates through the ejecta, a mildly relativistic cocoon would be formed and the interaction between the cocoon and the ambient medium would accelerate electrons via external shock in a wide angle, so that the merger-nova photons (i.e., thermal emission from the ejecta) would be scattered into higher frequency via an inverse Compton (IC) process when they propagate through the cocoon shocked region. We find that the IC scattered component peaks at the X-ray band and it will reach its peak luminosity on the order of days (simultaneously with the merger-nova emission). With current X-ray detectors, such a late X-ray component could be detected out to 200 Mpc, depending on the merger remnant properties. It could serve as an important electromagnetic counterpart of gravitational-wave signals from NS–NS/BH mergers. Nevertheless, simultaneous detection of such a late X-ray signal and the merger-nova signal could shed light on the cocoon properties and the concrete structure of the jet.

Key words: gamma-ray burst; general – gravitational waves – radiation mechanisms: non-thermal

1. Introduction

The LIGO-Virgo Collaboration has reported five gravitational-wave (GW) detections (GW150914, GW151226, GW170104, GW170608, and GW170814) from black hole–black hole (BH–BH) mergers, and one GW detection (GW170817) from a neutron star–neutron star (NS–NS) merger, which means a new era of GW astronomy has arrived (Abbott et al. 2016a, 2016b, 2017a, 2017b, 2017c, 2017d). Nevertheless, the electromagnetic (EM) counterparts of GW170817 have also been detected, including a weak short-duration gamma-ray burst (sGRB; Goldstein et al. 2017), an optical/IR transient in the galaxy NGC 4993 (Coulter et al. 2017), a radio counterpart (Hallinan et al. 2017), and a late re-brightening X-ray signal (Troja et al. 2017). A large number of teams across the world contributed to the observation of this source using ground- and space-based telescopes (see Abbott et al. 2017e for details).

Before the detection of GW170817, many associated EM counterparts had already been proposed for NS–NS/BH mergers, and their relative brightness is essentially determined by the properties of the post-merger central remnant object. In general, the remnant for an NS–BH merger would be a BH. But an NS–NS merger could lead to a BH or a supramassive NS, depending on the total mass of the NS–NS system and the NS equation of state (Dai et al. 2006; Gao & Fan 2006; Zhang 2013; Lasky et al. 2014; Gao et al. 2016). It is generally believed that, for both cases, sGRBs and their afterglow emission are expected as one of the major EM counterparts of NS–NS/BH mergers (Eichler et al. 1989; Narayan et al. 1992; Berger 2014). On the other hand, for both cases, the merger remnant would be surrounded by a mildly isotropic, sub-relativistic ejecta (which is composed of the tidally ripped and dynamically launched materials during the merger and the matter launched from the neutrino-driven wind from the accretion disk or neutron star surface; Rezzolla et al. 2011; Bauswein et al. 2013; Hotokezaka et al. 2013; Lei et al. 2013; Rosswog et al. 2013; Fernández et al. 2015; Song & Liu 2017). These ejecta are mostly composed of neutron-rich materials, and the radioactivity of these materials and the decay of r-process nuclei would heat the ejecta and then power an optical/IR transient (Li & Paczyński 1998; Metzger et al. 2010). When the merger remnant is a BH, r-process related radioactivity would be the only heating source, so that the luminosity of the optical/IR transient would be roughly 10^4 times that of the nova luminosity (Metzger et al. 2010). But if the merger remnant is a supramassive NS, its magnetic dipole radiation could serve as an additional heating source to the ejecta (Gao et al. 2013; Zhang 2013), which could easily exceed the r-process power, so that the thermal emission from the ejecta would be significantly enhanced (Yu et al. 2013; Metzger & Piro 2014). The luminosity of the optical transient, in this case, would be systematically brighter by more than one order of magnitude than the r-process dominating cases (Gao et al. 2017). Since the thermal emissions from the ejecta are essentially isotropic and also non-relativistic or mildly relativistic (due to the heavy mass loading), they therefore can be detected from any direction if the flux is high enough (see Metzger 2017a for a review).

All observed GW170817 EM counterparts are predicted, but some of them show unexpected behavior (see Metzger 2017b for a review). For instance, the fluence (∼1.4 × 10^{-7} erg cm^{-2}) of the associated sGRB 170817A falls into the average portion of the distributions of known sGRBs, but its peak isotropic luminosity (∼1.7 × 10^{47} erg s^{-1}) is abnormally low compared with other sGRBs (Goldstein et al. 2017). Considering the relatively large upper limit of the binary inclination angle relative to our line of sight from the GW signal analysis (Abbott et al. 2017d) and the self-consistency between γ-ray, X-ray, optical/IR, and radio observations, a cocoon emission has been proposed as one of the most...
concordant models (Gottlieb et al. 2017; Kasliwal et al. 2017; Piro & Kollmeier 2017; Xiao et al. 2017). Specifically, when a relativistic jet propagates through the surrounding ejecta, a mildly relativistic cocoon is formed embracing the main jet, so that some relatively weak emission would be expected from a wide-angle structure located in the periphery of the jet to explain the observed $\gamma$-ray emission (Gottlieb et al. 2018; Kathirgamaraju et al. 2018; Lazzati et al. 2017a, 2017b).

Within this scenario, the X-ray and radio observations could be interpreted as the off-axis afterglow from the main jet or from the on-axis afterglow emission of the cocoon, depending on the concrete properties of the jet and the ejecta (Gottlieb et al. 2017; Guidorzi et al. 2017; Kasliwal et al. 2017; Piro & Kollmeier 2017; Troja et al. 2017).

If this interpretation is correct, probably all NS–NS mergers would generate a wide-angle mildly relativistic cocoon. The interaction between the relativistic cocoon and the ambient medium could generate an external shock, where particles are believed to be accelerated, giving rise to broadband synchrotron radiation (cocoon afterglow emission). Considering the large opening angle of the cocoon, photons from other isotropic emission components, such as the thermal component from the ejecta, would be scattered into higher frequency via an inverse Compton (IC) process when they propagate through the cocoon external shock region. In this work, we will estimate the peak frequency and peak flux for this new emission component and will show the dependability of such emission by the currently available X-ray telescopes.

2. Seed Photons from Ejecta Thermal Emission

2.1. Numerical Model

Considering that one NS–NS merger event leaves behind a central remnant (either BH or NS), surrounded by a neutron-rich ejecta with mass $M_{ej}$ and initial dimensionless speed $\beta$. In any case, the ejecta would receive heating from the radioactive decay of the heavy nuclei synthesized in the ejecta via the $r$-process. If the central remnant is an NS, its magnetic dipole radiation could serve as an additional heating source to the ejecta. Moreover, the NS wind would continuously push from behind and accelerate the ejecta. Considering the energy dissipation through sweeping up the ambient medium, the dynamical evolution of the ejecta can be determined by (Yu et al. 2013)

$$\frac{d\Gamma}{dt} = \frac{dE}{dt} + \Gamma D \left( \frac{dM_{ej}}{dt} \right) _e - (\Gamma^2 - 1) c^2 \left( \frac{dM_{sw}}{dt} \right) ,$$

where $\Gamma$ is the bulk Lorentz factor, $t$ is the observer’s time in the jet frame, $D = 1/\Gamma(1 - \beta)$ is the Doppler factor, $E_{int}$ is the internal energy in the comoving frame, $t'$ is the time in the comoving frame, and $M_{sw} = -\frac{2}{3} R n_{mp} t'$ is the shock swept mass from the interstellar medium (with density $n$, where $R$ is the radius of the ejecta in the lab frame. With energy conservation, we have

$$\frac{dE}{dt} = L_{cen} + D^2 L_{ra}' - D^2 L_{e}' .$$

where $L_{cen}$ is the injected energy from the central engine, $L_{ra}'$ is the comoving radioactive power, and $L_{e}'$ is the comoving radiated bolometric luminosity. When the central engine is a BH, we normally expect $L_{cen} = 0$ (see Ma et al. 2017 for another opinion). When the central engine is an NS, normally a fraction $\xi$ of the dipole radiation luminosity is assumed to be injected into the ejecta (see Yu et al. 2013 for details), i.e., $L_{cen} = \xi L_{d}$, where $L_{d} = L_{sd} \left( 1 + \frac{1}{\Gamma^2} \right) ^{-2}$, with $L_{sd} = 10^{37} R_{s}^6 B_{i}^2 P_{e}^{-4} \text{erg s}^{-1}$ being the spin-down luminosity and $t_{sd} = 2 \times 10^3 R_{s}^6 B_{i}^2 P_{e}^{-4} \text{t}_{-3}^{-3}$ s being the spin-down timescale, where $P_{i}$, $B$, and $R_{s}$ are the initial spin period, the dipole magnetic strength, and the radius of the NS. Throughout this Letter, the convention $Q = 10^{6} Q_{6}$ is used for cgs units, except for the ejecta mass $M_{ej}$, which is in units of solar mass $M_{\odot}$.

Here, we adopt the empirical expression for the comoving radioactive power proposed by Korobkin et al. (2012):

$$L_{ra} = 4 \times 10^{39} M_{ej}, \left[ \frac{1}{2} - \frac{1}{\pi} \arctan \left( \frac{t'}{t_{\sigma}} \right) \right] ^{1.3} \text{erg s}^{-1},$$

where $t_{\sigma} \sim 1.3$ s and $t_{\rho} \sim 0.11$ s. The radiated bolometric luminosity could be expressed as

$$L_e = \begin{cases} \frac{E_{int} c}{\tau R / \Gamma}, & \tau > 1, \\ \frac{E_{int} c}{R / \Gamma}, & \tau < 1, \end{cases}$$

where $\tau = \kappa (M_{ej}/V') (R / \Gamma)$ is the optical depth of the ejecta with $\kappa$ being the opacity (Kasen & Bildsten 2010; Kotera et al. 2013).

The variation of the comoving internal energy $E_{int}$ could be expressed as (e.g., Kasen & Bildsten 2010; Yu et al. 2013)

$$\frac{dE_{int}}{dt'} = \xi D^{-2} L_{d} + L_{ra}' - L_{e}' - \frac{P'}{4} \frac{dV'}{dt'},$$

where $P' = E_{int}' / 3 V'$ is the radiation dominated pressure. The comoving volume evolution can be fully addressed by $dV'/dt' = 4 \pi R^2 \beta c$ together with $dR/dt = \beta c/(1 - \beta)$.

With $E_{int}$ being solvable, one can easily estimate the evolution of the effective temperature of the ejecta in the comoving frame $T_{eff}' = (E_{int}' / a' V' \max(\tau, 1))^{1/4}$. If a blackbody spectrum is assumed for the thermal emission from the ejecta, one can calculate the observed flux for a given frequency $\nu$:

$$F_{\nu} = \frac{1}{4} \frac{8 \pi \nu^2 D_l^2 \lambda^2}{h^2 c^2 \nu} \exp \left( \frac{h \nu}{D_l k T_{eff}'} - 1 \right),$$

where $k$ is the Boltzmann constant, $a'$ is the radiation constant, $h$ is the Planck constant, and $D_l$ is the luminosity distance.

2.2. Analytical Estimation

In principle, one can apply the above numerical modeling to calculate the observed flux of the thermal emission from the ejecta (henceforth we call it merger-nova emission) at any time and any frequency. The merger-nova photons could serve as the seed photons to be scattered by the cocoon-medium shock into X-ray band, which will be discussed later in detail. In order to better understand the features of these seed photons, we present some analytic approximations for the results of the
merger-nova emission, such as the peaking time of the merger-nova, the peak frequency, and peak flux at that time.

2.2.1. Black Hole as Merger Remnant

When the merger remnant is a BH, following the analytical estimation from Metzger et al. (2010), the bolometric luminosity of the r-process-powered merger-nova would reach its peak when the photon diffusion timescale equals the expansion timescale, where the radius of ejecta is

\[
R_p = \left( \frac{B v_e M_{ej}}{c} \right)^{1/2} \approx 3.74 \times 10^{14} \text{ cm} \left( \frac{v_e}{0.1c} \right)^{1/2} \frac{M_{ej}}{10^{-2}M_\odot}^{1/2} \left( \frac{\kappa}{1 \text{ cm}^2 \text{ g}^{-1}} \right)^{1/2},
\]

where \( M_{ej} \), \( v_e \), and \( \kappa \) are the mass, velocity, and opacity of the ejecta; and \( B \approx 0.07 \) for spherical outflow (e.g., Padmanabhan 2013). Under the assumption of free expansion, \( R_p \) would be reached on a timescale

\[
t_p = 1.25 \times 10^5 \text{ s} \left( \frac{v_e}{0.1c} \right)^{-1/2} \left( \frac{M_{ej}}{10^{-2}M_\odot} \right)^{1/2} \left( \frac{\kappa}{1 \text{ cm}^2 \text{ g}^{-1}} \right)^{-1/2}.
\]

The peak bolometric luminosity is given as

\[
L_p = 1.44 \times 10^{41} \text{ erg s}^{-1} \left( \frac{f}{10^{-6}} \right) \left( \frac{v_e}{0.1c} \right)^{1/2} \frac{M_{ej}}{10^{-2}M_\odot}^{1/2} \left( \frac{\kappa}{1 \text{ cm}^2 \text{ g}^{-1}} \right)^{-1/2},
\]

where \( f \) is a dimensionless number characterizing the heating generation rate (Li & Paczyński 1998; Metzger et al. 2010).

Using the Stefan–Boltzmann law, the effective temperature is roughly estimated as

\[
T_p = 6.17 \times 10^3 \text{ K} \left( \frac{f}{10^{-6}} \right)^{1/4} \left( \frac{v_e}{0.1c} \right)^{-1/8} \frac{M_{ej}}{10^{-2}M_\odot}^{-1/8} \left( \frac{\kappa}{1 \text{ cm}^2 \text{ g}^{-1}} \right)^{-3/8}.
\]

Assuming a blackbody spectrum, we can estimate the peak frequency of the thermal emission as

\[
\nu_p = 3.63 \times 10^{14} \text{ Hz} \left( \frac{f}{10^{-6}} \right)^{1/4} \left( \frac{v_e}{0.1c} \right)^{-1/8} \frac{M_{ej}}{10^{-2}M_\odot}^{-1/8} \left( \frac{\kappa}{1 \text{ cm}^2 \text{ g}^{-1}} \right)^{-3/8}.
\]

The according peak flux could be given as

\[
F_p = 3.15 \times 10^2 \text{ \mu Jy} \left( \frac{f}{10^{-6}} \right)^{3/4} \left( \frac{v_e}{0.1c} \right)^{5/8} \frac{M_{ej}}{10^{-2}M_\odot}^{5/8} \left( \frac{d}{10^{26}} \right)^{-2} \left( \frac{\kappa}{1 \text{ cm}^2 \text{ g}^{-1}} \right)^{-1/8}.
\]

2.2.2. Massive Neutron Star as a Merger Remnant

If the equation of state of nuclear matter is stiff enough, the central product for a binary neutron star merger could be a stable or a supramassive NS rather than a black hole. This newborn massive NS would be rotating with a rotation period on the order of milliseconds and may also contain a strong magnetic field \( B \geq 10^{14} \text{ G} \) similar to “magnetars.” In this case, the magnetar dipole radiation could easily dominate the heating and accelerating process for the ejecta so that the dynamics of the ejecta could be defined by energy conservation (see Gao et al. 2013 for details):

\[
\xi L dt = (\Gamma - 1)M_{ej}c^2 + (\Gamma^2 - 1)M_{tot}c^2.
\]

Considering that the number density of ambient medium in the NS–NS merger scenario is usually low, for most situations, the deceleration time for the ejecta should be larger than the spin-down timescale of the NS. At this stage, we have \((\Gamma - 1)M_{ej}c^2 \gg (\Gamma^2 - 1)M_{tot}c^2\), so that \(\Gamma = 1 \pm t\). With the energy injection from the millisecond magnetar, the ejecta could be accelerated into a mildly relativistic or even relativistic speed (Gao et al. 2013). In this case, we have the approximation of \(\Gamma \propto t\). With \(dR/dt = (3c)/(1 - \beta)\), we have \(R \propto t^2\).

The merger-nova emission peaks at \(t_p\) when \(\tau \approx 1\), where \(\tau = (\kappa M_{ej}/V)R\) is defined as the optical depth of the ejecta, \(\Gamma\) is the bulk Lorentz factor of the ejecta, and \(R\) and \(V\) are the radius and volume of the ejecta in the lab frame. The radius of the ejecta at \(t_p\) could be simply estimated as

\[
R_p \approx 6.92 \times 10^{14} \text{ cm} \kappa^{1/2} M_{ej}^{1/2}.
\]

When \(t_{sd} > t_p\), we can estimate the peak time in the observer frame as

\[
t_p = \left( \frac{R_p}{R_{sd}} \right)^{1/3} t_{sd}
\]

\[
\approx 1.54 \times 10^4 \kappa^{1/6} \xi^{-2/3} M_{ej,-3}^{5/6} B_{p,14}^{-4/3} P_{0,-3}^{8/3} R_{sd,4}^{-3/4},
\]

where \(R_{sd} \sim 2 \Gamma_{sd} \xi t_{sd}\) is the radius of the ejecta at the spin-down timescale, \(\Gamma_{sd} \approx \xi E_{rot}/M_{ej}c^2\) is the Lorentz factor of the ejecta at the spin-down timescale, and \(E_{rot} = (1/2)I_{sd} \Omega_{sd}^2 \approx 2 \times 10^{32} I_{sd} P_{0,-3}^2 \text{ erg}\) (with \(I_{sd} \sim 1.5\) for a massive neutron star) is the total spin energy of the millisecond magnetar.

At \(t_p\), the effective temperature of the ejecta in the observer frame is

\[
T_p = \left( \frac{\xi L_{sd}}{4\pi R_p^2 \sigma} \right)^{1/4}
\]

\[
\approx 1.31 \times 10^5 \kappa^{-1/4} \xi^{-1/4} M_{ej,-3}^{-1/4} B_{p,14}^{1/2} P_{0,-3}^{-1} R_{sd,4}^{3/2},
\]

where we assume a constant fraction (\(\xi\)) of the NS dipole luminosity to be thermalized in the outflow. With the numerical modeling, it is found that around the transparent time (i.e., \(\tau = 1\)), the radiated bolometric luminosity of the merger-nova is close to the injection energy power \(\xi L_{sd}\) (Yu et al. 2013).
Assuming a blackbody spectrum, we can estimate the peak frequency of the thermal emission in the observer frame as

\[ \nu_p \sim 7.70 \times 10^{15} \, \text{Hz} \, \kappa^{-1/4} \xi^{3/4} M_{1/4}^{-1/4} B_{p,\text{e}M}^{1/2} p_{2,5}^{-1/3} R_{3/2}^{3/2}. \]  

(17)

The according peak flux could be given by

\[ F_{\nu_p} \sim 1.03 \times 10^7 \, \mu Jy \, k^{1/4} \xi^{3/4} M_{1/4}^{3/4} R_{3/2}^{-1/2} d_{26}^{-2}. \]  

(18)

Similarly, for \( t_{\text{ad}} < t_p \) case, we have \( \Gamma \propto t \) and \( R \propto t^3 \) when \( t < t_{\text{ad}} \) and \( \Gamma \propto t^0 \) and \( R \propto t \) when \( t > t_{\text{ad}} \) (Gao et al., 2013), where \( L_{\text{ad}} \approx L_{\text{ad}}(t/t_{\text{ad}})^{-2} \) is taken when \( t > t_{\text{ad}} \). Then we can estimate the peak time of merger-nova as

\[ t_p \sim 4.07 \times 10^3 \, s \, k^{1/2} \xi^{-2} M_{1/3}^{2/3} p_{2,5}^{-2}. \]  

(19)

In this case, the temperature of the ejecta at \( t_p \) is

\[ T_p \sim 2.00 \times 10^5 \, K \, k^{-1/2} \xi^{5/2} M_{1/3}^{-3/2} B_{p,15}^{-1/2} p_{2,5}^{-3/2} R_{3/2}^{-3/2}. \]  

(20)

The peak frequency of the merger-nova emission in the observer frame is

\[ \nu_p \sim 1.18 \times 10^{16} \, \text{Hz} \, k^{-1/2} \xi^{5/2} M_{1/3}^{-3/2} B_{p,15,5}^{1/2} p_{2,5}^{-2} R_{3/2}^{-3/2}, \]  

(21)

and the corresponding peak flux is

\[ F_{\nu_p} \sim 3.66 \times 10^7 \, \mu Jy \, k^{-1/2} \xi^{15/4} M_{1/3}^{-7/2} B_{p,15,5}^{1/2} R_{3/2}^{-3/2} \times P_{\text{eM}}^{-6} d_{26}^{-2}. \]  

(22)

### 2.3. Inverse Compton Scattering

During the propagation of the cocoon, an external shock would form upon interaction with the ambient medium. The shock-accelerated electrons behind the blast wave are usually assumed to be distributed with a power-law function of electron energy, with a minimum Lorentz factor \( \gamma_m \): \( N(\gamma) d\gamma \propto \gamma_m^{-\gamma} d\gamma \), \( \gamma \geq \gamma_m \). Assuming that a constant fraction \( \epsilon_e \) of the shock energy is distributed to electrons, the minimum injected electron Lorentz factor can be estimated as

\[ \gamma_m = \epsilon_e \left( \frac{p - 2}{p - 1} \right) \frac{m_p}{m_e} \Gamma_{\text{co}} \approx 42.2 \epsilon_e^{-2} \Gamma_{\text{co}}, \]  

(23)

where \( \Gamma_{\text{co}} \) is the Lorentz factor of the cocoon, and \( p = 2.3 \) is adopted as commonly used in GRB afterglow modeling (Kumar & Zhang, 2015).

When the seed photons (with frequency \( \nu \)) from the merger-nova propagate through the cocoon external shock region, they would be scattered into higher frequency via the IC process. The typical photon frequency of the IC scattered merger-nova would be estimated as \(^1\) (Sari & Esin, 2001)

\[ \nu_{\text{IC}} \sim 2. \gamma_m^2 \nu_p. \]  

(24)

When merger remnant is a BH, we have

\[ \nu_{\text{IC, BH}} \sim 1.29 \times 10^{18} \, \text{Hz} \, \epsilon_e^{-2} \left( \frac{f}{10^{-6}} \right)^{1/4} \left( \frac{\nu_e}{0.1c} \right)^{-1/8} \times \left( \frac{M_{\text{ej}}}{10^{-2} M_{\odot}} \right)^{-1/8} \left( \frac{\kappa}{1 \, \text{cm}^2 \, \text{g}^{-1}} \right)^{-3/8}. \]  

(25)

When the merger remnant is an NS, we have

\[ \nu_{\text{IC, NS}} \sim 2.74 \times 10^{19} \, \text{Hz} \, \epsilon_e^{-2} \kappa^{-1/4} \xi^{1/4} \epsilon_{B,15,5}^{-1/2} B_{p,15,5}^{1/2} p_{2,5}^{-1} R_{3/2}^{3/2} \]  

(26)

where \( t_{\text{ad}} > t_p \), and \( t_{\text{ad}} < t_p \).

The peak flux for the IC scattering component could be estimated as

\[ F_{\nu_{\text{IC}}} \sim \tau_{\text{IC}} F_{\nu_p}, \]  

(28)

where \( \tau_{\text{IC}} \) is the optical depth for IC scatterings for a constant density medium (Sari & Esin, 2001), and \( \sigma_T \) is a Thompson scattering cross section. We thus have

\[ F_{\nu_{\text{IC,BH}}} \sim 2.62 \times 10^{-5} \, \mu Jy \left( \frac{d}{10^{16}} \right)^{-3/4} \left( \frac{\nu_e}{0.1c} \right)^{1/8} \times \left( \frac{M_{\text{ej}}}{10^{-2} M_{\odot}} \right)^{9/8} \left( \frac{d}{10^{16}} \right)^{-2} \times \left( \frac{\kappa}{1 \, \text{cm}^2 \, \text{g}^{-1}} \right)^{-3/8} \left( \frac{n}{1 \, \text{cm}^{-3}} \right)^{3/8}. \]  

(29)

\(^1\) In principle, one needs to first transform the seed photon energy from the observer frame to the cocoon-medium shock front frame, then calculate the inverse Compton scattering in that frame, and finally transform the scattered photon energy back to the observer frame. For simplicity, we assume that the seed photon injection direction is in line with the moving direction of the cocoon-medium shock front, in which case the two-step relativistic transformations could be canceled out.
and
\[ F_{\text{IC}, \text{NS}}^{\text{IC}} \sim 2.11 \times 10^{-1} \mu \text{Jy} \, \kappa^{5/12} \xi^{1/12} M_{13/12}^{1/12} \times B_{p}^{1/6} P_{-5/3}^{1/6} R_{6}^{-2/3} n_{d}^{-2/3}, \]

and
\[ F_{\text{IC}, \text{NS}}^{\text{IC}} \sim 1.98 \times 10^{-1} \mu \text{Jy} \, \xi^{7/4} M_{1-3}^{-1} \times B_{p}^{-1/2} P_{-2}^{-2} R_{6}^{-9/2} n_{d}^{-2}, \]

for cases \( t_{\text{ad}} > t_{p} \) and \( t_{\text{ad}} < t_{p} \), respectively.

3. Detectability

We use the numerical method described in Section 2.1 to calculate the light curve of the inverse Compton scattered merger-nova emission. Given some fiducial parameters, such as the ejecta mass, velocity, and opacity as \( 10^{-2} M_{\odot}, 0.2 c, \) and 1 cm\(^{-2}\)g\(^{-1}\), and the magnetar parameters as \( B = 10^{14} \text{G}, P_{0} = 1 \text{ms}, R_{\text{ej}} = 10^{6} \text{cm} \), and \( \xi = 0.3 \), we find that the IC scattered merger-nova peaks at the X-ray band. In Figure 1, we compare the peak flux of the IC scattered merger-nova with the sensitivity of current X-ray facilities, such as Swift, XRT, Chandra, and XMM-Newton (Burrows et al. 2005; Jansen et al. 2001; Weisskopf et al. 2002). We find that for NS–BH mergers or NS–NS mergers with BH being the merger remnant, such an X-ray component is only detectable out to 2 Mpc with currently available facilities. However, for NS–NS mergers with a massive NS being the merger remnant, the IC scattered X-ray component would become much brighter, and it will be detectable out to 200 Mpc (the designed horizon of aLIGO for NS–NS mergers; Abbott et al. 2009). In Figure 1, we also plot the expected IC scattered merger-nova for GW170817, and we find that it is too dim to account for the late Chandra X-ray observations.

4. Conclusion and Discussion

The recent observations of GW170817 and its EM counterparts have proven the prediction that NS–NS/BH mergers could lead to rich and bright EM emissions, invoking several emission components. In general, the jet component would give rise to an sGRB and its afterglow, and the isotropic component would give rise to a merger-nova emission and its afterglow. Recent studies suggest that the jet component is structured, with a relativistic jet surrounded by a mild relativistic cocoon. In this case, the interaction between the cocoon and the ambient medium would accelerate electrons via external shock in a wide angle, so that the merger-nova photons would be scattered into higher frequency via the IC process when they propagate through the cocoon external shock region.

In this work, we find that the IC scattered component peaks at the X-ray band and it will reach its peak luminosity simultaneously with the merger-nova. For NS–BH mergers or NS–NS mergers with BH being the merger remnant, the X-ray component is detectable out to 2 Mpc with current facilities (such as Chandra and XMM-Newton). On the other hand, if the total mass of binary neutron star system is small enough and the equation of state of nuclear matter is stiff enough, the merger of two NSs could leave behind a supramassive NS. In this case, the merger-nova emission could be significantly enhanced so that the IC scattered X-ray component also becomes brighter. It will be detectable out to 200 Mpc with current facilities. Note that even for a BH remnant case, the magnetic wind driven by the Blandford–Payne process (Blandford & Payne 1982) from a newborn BH accretion disk or fallback accretion disk would significantly enhance the merger-nova emission (see Chen et al. 2017; Ma et al. 2017 for details); in this case, the IC scattered X-ray component could become as bright as the magnetar remnant case.

Our newly proposed late X-ray emission could serve as an important EM counterpart of GW signals. Simultaneous detection of such an X-ray signal and the merger-nova signal could help to investigate the cocoon properties and the concrete structure of the jet.

It is worth noting that some other mechanism could also generate late X-ray emission, which may outshine the proposed signal here. For instance, if the merger remnant of NS–NS is a supramassive NS, the X-rays powered by NS wind dissipation would diffuse out at late time when the ejecta becomes (or be close to) optically thin, a late X-ray re-brightening would be expected (Metzger & Piro 2014; Gao et al. 2015, 2017). But if the supramassive NS has collapsed into a black hole before the surrounding ejecta become transparent, such a signal would disappear. On the other hand, the afterglow emission from the structure jet could also provide X-ray photons, but its strength sensitively depends on the viewing angle and the energy distribution within the jet. For most proper viewing angles, the corresponding jet energy is usually small so that a relatively weak X-ray afterglow emission is expected (Lazzati et al. 2017b).

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