BRIGHT BROADBAND AFTERGLOWS OF GRAVITATIONAL WAVE BURSTS FROM MERGERS OF BINARY NEUTRON STARS

He Gao1,2, Xuan Ding1, Xue-Feng Wu1,3,4, Bing Zhang2,5,6, and Zi-Gao Dai7

1 Purple Mountain Observatory, Chinese Academy of Sciences, Nanjing 210008, China; xfwu@pmo.ac.cn
2 Department of Physics and Astronomy, University of Nevada, Las Vegas, NV 89154, USA; zhang@physics.unlv.edu
3 Chinese Center for Antarctic Astronomy, Chinese Academy of Sciences, Nanjing 210008, China
4 Joint Center for Particle Nuclear Physics and Cosmology of Purple Mountain Observatory–Nanjing University, Chinese Academy of Sciences, Nanjing 210008, China
5 Department of Astronomy, Peking University, Beijing 100871, China
6 Kavli Institute of Astronomy and Astrophysics, Peking University, Beijing 100871, China
7 School of Astronomy and Space Science, Nanjing University, Nanjing 210093, China; dzz@nju.edu.cn

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ABSTRACT

If double neutron star mergers leave behind a massive magnetar rather than a black hole, then a bright early afterglow can follow the gravitational wave burst (GWB) even if there is no short gamma-ray burst (SGRB)–GWB association or if there is an association but the SGRB does not beam toward Earth. Besides directly dissipating the proto-magnetar wind, as suggested by Zhang, here we suggest that the magnetar wind could push the ejecta launched during the merger process and, under certain conditions, would reach a relativistic speed. Such a magnetar-powered ejecta, when interacting with the ambient medium, would develop a bright broadband afterglow due to synchrotron radiation. We study this physical scenario in detail and present the predicted X-ray, optical, and radio light curves. We show that the X-ray and optical light curves usually peak around the magnetar spin-down timescale ($\sim 10^3$–$10^5$ s), reaching brightnesses readily detectable by wide-field X-ray and optical telescopes, and remain detectable for an extended period. The radio afterglow peaks later, but is much brighter than the case without a magnetar energy injection. Therefore, such bright broadband afterglows, if detected and combined with GWBs in the future, would be a probe of massive millisecond magnetars and stiff equations of state for nuclear matter.

Key words: gravitational waves – hydrodynamics – radiation mechanisms: non-thermal – shock waves – stars: magnetars – stars: neutron

Online-only material: color figures

1. INTRODUCTION

The next generation of gravitational wave (GW) detectors, such as the Advanced LIGO (Abbott et al. 2009), Advanced VIRGO (Acernese et al. 2008), and KAGRA (Kuroda et al. 2010) interferometers, are expected to detect GW signals from mergers of two compact objects. These gravitational wave bursts (GWBs) have a well-defined “chirp” signal, which can be unambiguously identified. Once detected, the GW signals would open a brand new channel for us to study the universe, especially the physics in the strong field regime. Due to the faint nature of GWs, an associated electromagnetic (EM) emission signal coinciding with a GWB in both trigger time and direction would increase the signal-to-noise ratio of the GW signal and therefore would be essential for its identification.

One of the top candidates of GWBs is the merger of two neutron stars (i.e., NS–NS mergers; Taylor & Weisberg 1982; Kramer et al. 2006). The EM signals associated with such an event include a short gamma-ray burst (SGRB; Eichler et al. 1989; Rosswog et al. 2013; Gehrels et al. 2005; Barthelmy et al. 2006; Aloy et al. 2005; Shibata et al. 2005; Rosswog et al. 2013; Hotokezaka et al. 2013). The radioactivity of this ejecta powers the macronova and the interaction between the ejecta and the ambient medium is the source of radio afterglow. Usually, the merger product is assumed to be a black hole or a temporal hyper-massive neutron star which survives 10–100 ms before collapsing into the black hole (e.g., Rosswog et al. 2003, 2013; Aloy et al. 2005; Shibata et al. 2005; Rezzolla et al. 2011). Nonetheless, recent observations of Galactic neutron stars and NS–NS binaries suggest that the maximum neutron star mass can be high, which is close to the total mass of the NS–NS systems (Dai et al. 2006; Zhang 2013 and references therein). Indeed, for the measured parameters of six known Galactic neutron star binaries and a range of equations of state, the majority of mergers of the known binaries will form a massive millisecond pulsar and survive for an extended period of time (Morrison et al. 2004). When the equation of state of nuclear matter is stiff (see arguments in Dai et al. 2006; Zhang 2013 and references therein), a stable massive neutron star would form after the merger. This newborn massive neutron star would be differentially rotating. The dynamo mechanism may operate and generate an ultrastrong magnetic field (Duncan & Thompson 1992; Kluzniak & Ruderman 1998; Dai & Lu 1998b), so that the product is very likely a millisecond magnetar. Evidence of a magnetar following some SGRBs has been collected in the Swift data (Rowlinson et al. 2010; Rowlinson et al. 2013), and magnetic activities of such a post-merger massive neutron star have been suggested to interpret several X-ray flares and plateau phase in SGRBs (Dai et al. 2006; Gao & Fan 2006; Fan & Xu 2006).
Since both the GW signal and the millisecond magnetar wind are nearly isotropic, a bright EM signal can be associated with an NS–NS merger GWB regardless of whether there is an SGRB–GWB association (Zhang 2013). Even if there is an association, most GWBs would not be associated with the SGRB since SGRBs are collimated. Zhang (2013) proposed that the near-isotropic magnetar wind of a post-merger millisecond magnetar would undergo magnetic dissipation (Zhang & Yan 2011) and power a bright X-ray afterglow emission. Here we suggest that after partially dissipating the magnetic energy, a significant fraction \( \xi \) of the magnetar spin energy would be used to push the ejecta, which drives a strong forward shock into the ambient medium. The continuous injection of the Poynting flux into the blast wave modifies the blast wave dynamics and leads to rich radiation signatures (Dai & Lu 1998a; Zhang & Mészáros 2001; Dai 2004). Figure 1 presents a physical picture of several EM emission components appearing after the merger. A massive millisecond magnetar is formed at the central engine. Near the spin axis, there might be an SGRB jet. An observer toward this jet (red observer) would see an SGRB. At even larger angles (orange observer), the magnetar wind is confined by the axis, there might be an SGRB jet. An observer toward this jet (red observer) would see an SGRB. At larger angles (yellow observer), a free magnetar wind may be released, whose dissipation would power a bright X-ray afterglow (Zhang 2013). At even larger angles (orange observer), the magnetar wind is confined by the axis, there might be an SGRB jet. An observer toward this jet (red observer) would see an SGRB. At larger angles (yellow observer), a free magnetar wind may be released, whose dissipation would power a bright X-ray afterglow (Zhang 2013). After releasing some dissipated energy, a significant fraction of the spinning energy would push the ejecta and shock into the ambient medium (Dai & Lu 1998a; Zhang & Mészáros 2001). Synchrotron emission from the shocked medium (red shell) would power brighter X-ray, optical, and radio afterglow emission, which is calculated in this work.

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

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2. THE MODEL

The post-merger hyper-massive neutron star may be near the breakup limit, so that the total spin energy \( E_{\text{rot}} = (1/2) I \Omega_0^2 \simeq 2 \times 10^{52} I_{45} P_{0,-3}^2 \text{erg} \) (with \( I_{45} \sim 1.5 \) for a massive neutron star) may be universal. Here \( P_0 \sim 1 \) ms is the initial spin period of the proto-magnetar. Throughout the paper, the convention \( Q = 10^9 Q_n \) is used for cgs units, except for the ejecta mass \( M_{ej} \), which is in units of solar mass \( M_\odot \). Given nearly the same total energy, the spin-down luminosity and the characteristic spin-down timescale critically depend on the polar-cap dipole magnetic field strength \( B_p \) (Zhang & Mészáros 2001), i.e., \( L_{sd} = L_{sd,0}/(1+t/T_{sd})^2 \), where \( L_{sd,0} \simeq 10^{49} \text{erg s}^{-1} B_p^{2.15} R_6^6 P_0^{-3} \), and the spin-down timescale \( T_{sd} \simeq 2 \times 10^3 \text{s} I_{45} B_p^{-2.15} P_0^{-3} R_6^{-6} \simeq E_{rot}/L_{sd,0} \), where \( R = 10^6 R_6 \text{cm} \) is the stellar radius.8

After the internal dissipation of the magnetar wind that powers the early X-ray afterglow (Zhang 2013), the remaining spin energy would be added to the blast wave. The dynamics of the blast wave depends on the magnetization parameter \( \sigma \) of the magnetar wind after the internal dissipation. Since for the confined wind magnetic dissipation occurs upon interaction between the wind and the ejecta, in this paper, we assume that the wind is still magnetized (moderately high \( \sigma \)), so that there is no strong reverse shock into the magnetar wind (Zhang & Kobayashi 2005; Mimica et al. 2009).9 As a result, the remaining spin energy is continuously injected into the blast wave with a luminosity \( L_0 = \xi L_{sd,0} \), where \( \xi < 1 \) denotes the fraction of the spin-down luminosity that is added to the blast wave. The evolution of the blast wave can be described by a system...
with continuous energy injection (Dai & Lu 1998a; Zhang & Mészáros 2001).

The newly formed massive magnetar is initially hot. A Poynting-flux-dominated outflow is launched ~10 s later, when the neutrino-driven wind is clean enough (Metzger et al. 2011).

At this time, the front of the ejecta traveled a distance of ~6 × 10^{10} cm (for u ∼ 0.2c), with a width ∆ ~ 10^{7} cm. The ultrarelativistic magnetar wind takes ~2 s to catch up the ejecta and drives a forward shock into the ejecta. Balancing the magnetic pressure and the ram pressure of shocked fluid in the ejecta, one can estimate the shocked fluid speed as \( v_s \sim 10^{-3} c T_4^{-1/2} M_0^{-1/2} \), which is in the same order of magnitude as the shock speed.

Thus, the forward shock would cross the ejecta, and the blast wave ploughs into the ambient medium. The dynamics of the blast wave during this stage is defined by energy conservation: \( E_{\text{tot}} = (\gamma - 1) M_{\text{ej}} c^2 + (\gamma^2 - 1) M_{\text{sw}} c^2 \),

\[ L_{\text{tot}} = (\gamma - 1) M_{\text{ej}} c^2 + (\gamma^2 - 1) M_{\text{sw}} c^2, \tag{1} \]

where \( M_{\text{sw}} = (4\pi/3)R^3 n_m\rho \) is the swept mass from the interstellar medium. Initially, \((\gamma - 1) M_{\text{ej}} c^2 \gg (\gamma^2 - 1) M_{\text{sw}} c^2\) with time until \( t = \text{min}(T_{\text{dec}}, T_{\text{sd}})\), where the deceleration timescale \( T_{\text{dec}} \) is defined by the condition \((\gamma - 1) M_{\text{ej}} c^2 = (\gamma^2 - 1) M_{\text{sw}} c^2\). By setting \( T_{\text{dec}} \sim T_{\text{sd}}\), we can derive a critical ejecta mass

\[ M_{\text{ej,c1}} \sim 10^{-3} M_0 n_4 L_4^{1/5} P_{-3}^{1/5} \frac{5}{4} T_{16}^{3/8} P_{-3}^{5/2} c, \]

\[ \sim 10^{-3} M_0 n_4 L_4^{1/5} P_{14}^{3/8} B_{-5}^{2/3} R_6^{-9/4} P_{-3}^{4/5} c^{7/8}, \tag{2} \]

which separates regimes with different blast wave dynamics. For a millisecond dynamic magnetar, the parameters \( L_4, R_6, n_4, P_{-3} \) are all essentially fixed values. The dependence on \( n \) is very weak (1/8 power), so the key parameters that determine the blast wave parameters are the ejecta mass \( M_{\text{ej}} \) and the magnetar injection luminosity \( L_0 \) (or the magnetic field strength \( B_p \)). If \( M_3 \gg M_{\text{ej,c1}} \) (or \( T_{\text{dec}} < T_{\text{sd}} \)), then the ejecta can be accelerated linearly until the deceleration radius, after which the blast wave decelerates, but still with continuous energy injection until \( T_{\text{sd}} \).

Conversely, in the opposite regime \( M_3 \gg M_{\text{ej,c1}} \) or \( T_{\text{sd}} < T_{\text{dec}} \), the blast wave is only accelerated to \( T_{\text{dec}} \), after which it coasts before decelerating at \( T_{\text{dec}} \). In the intermediate regime of \( M_3 \sim M_{\text{ej,c1}} \) (or \( T_{\text{dec}} \sim T_{\text{sd}} \)), the blast wave decays after being linearly accelerated.

There is another critical ejecta mass which defines whether the blast wave can reach a relativistic speed. This is defined by \( E_{\text{tot}} = 2(\gamma - 1) M_{\text{ej}} c^2\). Defining a relativistic ejecta as \( \gamma - 1 > 1 \), this second critical ejecta mass is

\[ M_{\text{ej,c2}} \sim 6 \times 10^{-3} M_0 I_{45} P_{-3}^{2} c^2. \tag{3} \]

An ejecta heavier than this would not be accelerated to a relativistic speed.

Below, we discuss four dynamical regimes.

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Table 1: Expression of the Lorentz Factor and Radius as a Function of Model Parameters in Different Temporal Regimes for All Dynamical Cases

| Case | \( \gamma \) | \( R \) |
|------|--------------|---------|
| Case I: \( M_3 \) < \( M_{\text{ej,c1}} \) or \( T_{\text{sd}} > T_{\text{dec}} \) | Realistic | \( \gamma \) |
| Case II: \( M_3 \sim M_{\text{ej,c1}} \) or \( T_{\text{dec}} \sim T_{\text{sd}} \) | Relativistic | \( R \) |
| Case III: \( M_3 \gg M_{\text{ej,c1}} \) or \( T_{\text{dec}} < T_{\text{sd}} \) | Small | \( \gamma \) |

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Case I: \( M_3 < M_{\text{ej,c1}} \) or \( T_{\text{sd}} > T_{\text{dec}} \). This requires both a small \( L_0 \) (or low \( B_p \)) and a small \( M_3 \). We take an example with \( L_0 \sim 10^{47} \) erg s^{-1} (\( B_p \sim 10^{14} \) G) and \( M_3 \sim 10^{-4} \) M⊙. To describe the dynamics in such a case, besides the spin-down timescale \( T_{\text{sd}} \), we need three more characteristic timescales and the Lorentz factor value at the deceleration time

\[ T_{\text{dec}} \sim 4.4 \times 10^{4} s L_{0.47}^{7/10} M_{4.5}^{1/3} n_{-11}^{-1/10} \]

\[ T_{\text{N1}} \sim 3.6 \times 10^{5} s L_{0.47}^{1/5} M_{4.5}^{2/5} \]

\[ T_{\text{N2}} \sim 4.5 \times 10^{7} s L_{0.47}^{1/5} M_{4.5}^{2/5} \]

\[ \gamma_{\text{dec}} \sim 12.2 L_{0.47}^{3/10} M_{4.5}^{1/3} n_{-11}^{-1/10} + 1, \tag{4} \]

where \( T_{\text{N1}} \) and \( T_{\text{N2}} \) are the two timescales when the blast wave passes the non-relativistic-to-relativistic transition line \( \gamma - 1 = 1 \) during the acceleration and deceleration phases. With these parameters, one can characterize the dynamical evolution of the blast wave (Figure 2(a)), as shown in Table 1. Based on the dynamics, we can quantify the temporal evolution of synchrotron radiation characteristic frequencies \( \nu_{\text{r}}, \nu_{\text{m}}, \) and \( \nu_{\text{c}} \), and the peak flux, \( F_{\text{r, max}} \). The evolutions of the characteristic frequencies are presented in Figure 2(b) and collected in Table 2.

Following the standard procedure in Sari et al. (1998), we derive the synchrotron radiation characteristic frequencies and the peak flux density at \( T_{\text{dec}} \),

\[ \nu_{\text{a,dec}} \sim 5.0 \times 10^{8} Hz L_{0.47}^{5/30} M_{4.5}^{1/25} n_{29/50}^{1/5} \epsilon_{c,-1}^{1/5} (p - 2) / (p - 1)^{3/5} f(p)^{3/5} \]

\[ \nu_{\text{m,dec}} \sim 1.3 \times 10^{13} Hz L_{0.47}^{6/5} M_{4.5}^{1/5} n_{10/10}^{1/10} \epsilon_{c,-1}^{1/2} (p - 2) / (p - 1)^{2} \]

\[ \nu_{\text{c,dec}} \sim 9.6 \times 10^{14} Hz L_{0.47}^{12/5} M_{4.5}^{1/5} n_{9/10}^{1/2} \epsilon_{c,-1}^{2} (p - 2) / (p - 1)^{2} \]

\[ F_{\text{r, max,dec}} \sim 1.7 \times 10^{5} \mu Jy L_{0.47}^{1/10} M_{4.5}^{1/5} n_{29/50}^{1/2} D_{27}^{2}, \tag{5} \]
Figure 2. Calculation results for Case I: $L_0 \sim 10^{47}$ erg s$^{-1}$ and $M_{ej} \sim 10^{-4} M_\odot$ (for all examples, we adopt $\xi = 0.5$, $p = 2.3$). (a) The dynamical evolution of the parameter ($\gamma - 1$); (b) temporal evolutions of the characteristic frequencies $v_a$, $v_m$, and $v_c$, and the peak flux density $F_{\nu, max}$; (c) analytical light curve in the $R$ band (blue) and the 10 GHz radio band (red); (d) analytical light curve in the X-ray band. The solid and dashed lines represent $n = 1$ cm$^{-3}$ and $n = 10^{-3}$ cm$^{-3}$, respectively. In panels (c) and (d), we mark the spectral and temporal indices for each segment of the light curves for $n = 1$ cm$^{-3}$. The main figures denote the time regimes when the light curves are detectable. The insets show the full light curves for completeness. Both X-ray and optical light curves reach their peaks around $10^9$ s and remain detectable in years. The radio light curve peaks around $10^7$ s, and lasts even longer. The peak flux for X-ray, optical, and radio could be as bright as $10^{11}$ erg s$^{-1}$ cm$^{-2}$, 10 mJy, and Jy, respectively.

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

where $f(p) = (\Gamma((3p + 2)/12)\Gamma((3p + 9)/12))/(\Gamma((3p + 19)/12)\Gamma((3p - 1)/12))$. With the temporal evolution power-law indices of these parameters (Table 2), one can calculate the X-ray, optical, and radio afterglow light curves. Note that there are two more temporal segments listed in Table 2, since $v_a$ crosses $v_m$ twice at

$$T_{max} \sim 1.4 \times 10^2 s L_{0,47}^{-5/4} M_{ej, -4}^{5/4} \xi_{-1}^{-5/4} \epsilon_{-1}^{-1/8} \left( \frac{p - 2}{p - 1} \right)^{-5/4} \times (p + 1)^{1/4} f(p)^{1/4},$$

$$T_{ma2} \sim 1.9 \times 10^8 s L_{0,47}^{1/5} n_{-2}^{-2/5} T_{sd, 5}^{1/5} \xi_{-1}^{-1/8} \epsilon_{-1}^{-1/8} \left( \frac{p - 2}{p - 1} \right)^{2} \times (p + 1)^{-2/5} f(p)^{-2/5},$$

respectively. We present the light curves in the X-ray (Figure 2(d)), optical, and radio (10 GHz) bands (Figure 2(c)).

The distance is taken as 300 Mpc, the detection horizon of the Advanced LIGO.

Case II: $M_{ej} \sim M_{ej,c, 1}$ or $T_{sd} \sim T_{dec}$. The dynamics and the expressions of the characteristic parameters become simpler:

$$T_{dec} \sim T_{sd}$$

$$T_{N1} \sim 12 s \xi^{-1} M_{ej, -4} T_{sd, 3}$$

$$T_{N2} \sim 1.3 \times 10^8 s \xi^{-8/3} M_{ej, -4}^{-8/3} T_{sd, 3}$$

$$\gamma_{sd} \sim 83.3 \xi^{-1} M_{ej, -4}^{-1} + 1.$$

The temporal indices of the evolutions of $v_a$, $v_m$, $v_c$, and $F_{\nu, max}$ are listed in Table 2, and the expressions of $\gamma$ and $R$ are shown in Table 1.

As examples, we consider $L_0 \sim 10^{49}$ erg s$^{-1}$ ($B_p \sim 10^{15}$ G) versus $M_{ej} \sim 10^{-4} M_\odot$, which satisfies $T_{sd} \sim T_{dec}$. 

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For this example, the dynamics and the expressions of the characteristic parameters become

\[ T_{\text{dec}} \sim 1.5 \times 10^4 \, \text{s} \, \xi^{-3/7} M_{\text{ej}}^{6/7} n_{-1/3} \]
\[ T_{\text{N1}} \sim 59.9 \, \xi^{-1} M_{\text{ej}}^{-3} T_{\text{sd,3}} \]
\[ T_{\text{N2}} \sim 2.7 \times 10^4 \, \xi^{-1/3} n_{-1/3} \]
\[ \gamma_{\text{sd}} \sim 16.7 \, \xi^{-1} M_{\text{ej}}+1 \]

and

\[ v_{\text{sa, sd}} \sim 1.6 \times 10^8 \, \text{Hz} \, \xi^{3/5} M_{\text{ej}}^{4/5} n_{-1/3}^{4/5} T_{\text{sd,3}}^{-1} \epsilon_{-1} B_{-2}^{-1/8} \left( \frac{p-2}{p-1} \right)^{1/5} \]
\[ \times (p+1)^{3/5} f(p)^{1/5} \]
\[ v_{\text{v, sd}} \sim 4.5 \times 10^{15} \, \text{Hz} \, \xi^{-4} M_{\text{ej}}^{-1} n_{-1/3}^{1/2} \epsilon_{-1} B_{-2}^{-1/2} \left( \frac{p-2}{p-1} \right)^2 \]
\[ \times (p+1)^{3/5} f(p)^{1/5} \]
\[ v_{\text{c, sd}} \sim 5.3 \times 10^{17} \, \text{Hz} \, \xi^{-4} M_{\text{ej}}^{-1} n_{-1/3}^{-3/2} T_{\text{sd,3}}^{-2} B_{-2}^{-1/2} \]
\[ F_{\text{v, max, sd}} \sim 6.5 \times 10^{-2} \, \mu\text{Jy} \, \xi^{3/4} M_{\text{ej}}^{7/4} n_{-1/3}^{3/4} T_{\text{sd,3}}^{5/4} \epsilon_{-1} B_{-2}^{-2} D_{27}^{-2} \]
\[ T_{\text{ma1}} \sim 1.0 \, \xi^{-5/4} M_{\text{ej}}^{-1} n_{-1/3}^{1/2} \epsilon_{-1} B_{-2}^{-1/2} \left( \frac{p-2}{p-1} \right)^{-5/4} \]
\[ \times (p+1)^{1/5} f(p)^{1/5} \]
\[ T_{\text{ma2}} \sim 9.9 \times 10^7 \, \xi^{-1/2} n_{-1/3}^{1/2} \epsilon_{-1} B_{-2}^{-1/2} \left( \frac{p-2}{p-1} \right)^2 \]
\[ \times (p+1)^{-2/5} f(p)^{-2/5} \].

The power-law indices of various parameters for this case are also collected in Table 2, and the dynamics, frequency evolutions, and light curves are presented in Figure 4.

Case IV: \( M_{\text{ej}} > M_{\text{ej,c}} \). In this case, the blast wave never reaches a relativistic speed. The dynamics is similar to Case III, with the coating regime in the non-relativistic phase. The dynamics for a non-relativistic ejecta and its radio afterglow emission have been discussed in Nakar & Piran (2011). Our Case IV resembles what is discussed in Nakar & Piran (2011), but the afterglow flux is much enhanced because of the larger total energy involved.

### 3. DETECTABILITY AND IMPLICATIONS

For all the cases, bright broadband EM afterglow emission signals are predicted. The light curves typically show a sharp rise around \( T_{\text{sd}} \), which coincides with the ending time of the X-ray afterglow signal discussed by Zhang (2013) due to internal dissipation of the magnetar wind. The X-ray afterglow luminosity predicted in our model is generally lower than that of the internal dissipation signal, but the optical and radio signals are much brighter. In some cases, the \( K \)-band magnitude can reach 11th at 300 Mpc, if \( M_{\text{ej}} \) is small enough (so that the blast wave has a high Lorentz factor) and the medium density is not too low. The duration of detectable optical emission ranges from \( 10^3 \) s to a year timescale. The radio afterglow can reach the Jy level for an extended period of time, with peak reached in the year timescale. These signals can be readily picked up by all-sky optical monitors and radio surveys. The X-ray afterglow can also be picked up by large field-of-view imaging telescopes such as the ISS-Lobster.

Since these signals originate from interaction between the magnetar wind and the ejecta in the equatorial directions, they are not supposed to be accompanied with SGRBs and some...
internal-dissipation X-ray afterglows (Zhang 2013) in the free wind zone. Due to a larger solid angle, the event rate for this geometry (orange observer in Figure 1) should be higher than the other two geometries (red and yellow observers in Figure 1). However, the brightness of the afterglow critically depends on the unknown parameters such as $M_{ej}$, $B_p$ (and hence $L_0$), and $n$. The event rate also crucially depends on the event rate of NS–NS mergers and the fraction of mergers that leave behind a massive magnetar rather than a black hole.

This afterglow signal is much stronger than the afterglow signal due to ejecta–medium interaction with a black hole as the post-merger product (Nakar & Piran 2011). The main reason is the much larger energy budget involved in the magnetar case. Since the relativistic phase can be achieved, both the X-ray and optical afterglows are detectable, which peak around the magnetar spin-down timescale ($10^3$–$10^5$ s). The radio peak is later, similar to the black hole case (Nakar & Piran 2011), but the radio afterglow flux is also much brighter (reaching Jy level) due to a much larger energy budget involved. The current event rate limit of $>350$ mJy radio transients in the minutes-to-days timescale at 1.4 GHz is $<6 \times 10^{-4}$ deg$^{-2}$ yr$^{-1}$ (Bower & Saul 2011), or $<20$ yr all sky. In view of the large uncertainties in the NS–NS merger rate and the fraction of millisecond magnetar as the post-merger product, our prediction is entirely consistent with this upper limit. Because of their brightness, these radio transients can be detected outside the Advanced LIGO horizon, which may account for some submJy radio transients discovered by the Very Large Array (Bower et al. 2007).

Recently, Kyutoku et al. (2012) proposed another possible EM counterpart of GWB with a wide solid angle. They did not invoke a long-lasting millisecond magnetar as the merger product, but speculated that during the merger process a breakout shock from the merging neutron matter would accelerate a small fraction of surface material, which reaches a relativistic speed. Such an outflow would also emit broadband synchrotron emission by shocking the surrounding medium. Within that scenario, the predicted peak flux is lower and the duration is shorter.

![Figure 3](image-url)
than the EM signals predicted in Zhang (2013) and this work due to a much lower energy carried by the outflow.

Detecting the GWB-associated bright signals as discussed in this paper would unambiguously confirm the astrophysical origin of GWBs. Equally important, it would suggest that NS–NS mergers leave behind a hyper-massive neutron star, which gives an important constraint on the neutron star equation of state. With the GWB data, one can infer the information of the two neutron stars involved in the merger. Modeling afterglow emission can give useful constraints on the ejected mass $M_{ej}$ and the properties of the post-merger compact objects. Therefore, a combination of GWB and afterglow information would shed light into the detailed merger physics, and in particular, provide a probe of massive millisecond magnetars and stiff equations of state for neutron matter.

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Figure 4. Calculation results for Case III: $L_X \sim 10^{50}$ erg s$^{-1}$ and $M_{ej} \sim 10^{-3} M_\odot$. Descriptions of panels are the same as in Figure 2. For $n = 1$ cm$^{-3}$, both the X-ray and optical light curves reach their peaks around $10^5$ s, and the radio light curve peaks around $10^6$ s. The peak flux for X-ray, optical, and radio is $10^{-10}$ erg s$^{-1}$ cm$^{-2}$, 10 mJy, and Jy, respectively. Taking $R$-band magnitudes 20 and 10, the optical and X-ray afterglows are $\sim 10^7$ s and $\sim 10^6$ s, respectively. The radio duration lasts even longer. If $n = 10^{-3}$ cm$^{-3}$, the optical signal just reaches an $R$-band magnitude of 20 around $\sim 10^8$ s, while the X-ray afterglow is detectable with a duration of $\sim 10^8$ s.

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