Outflow-driven Transients from the Birth of Binary Black Holes. I. Tidally Locked Secondary Supernovae

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Abstract

We propose a new type of electromagnetic transient associated with the birth of binary black holes (BBHs), which may lead to merger events accompanied by gravitational waves in ~0.1–1 Gyr. We consider the newborn BBHs formed through the evolution of isolated massive stellar binaries. For a close massive binary, consisting of a primary black hole (BH) and a secondary Wolf–Rayet (WR) star that are orbiting around each other, the spin period of the WR star can be tidally locked or synchronized to its orbital period. Then the angular momentum of the outer material of the WR star is large enough to form an accretion disk around a newborn, secondary BH, following its core-collapse. This disk may produce an energetic outflow with a kinetic energy of ~10^{50}–10^{52} erg and an outflow velocity of ~10^{10} cm s^{-1}, resulting in an optical transient with an absolute magnitude from approximately −14 to approximately −17 with a duration of around a day. This type of transient also produces detectable radio signals ~1–10 years after the birth of BBHs, via synchrotron emission from nonthermal electrons accelerated at external shocks. The predicted optical transients have a shorter duration than ordinary core-collapse supernovae. Dedicated optical transient surveys can detect them and distinguish them from ordinary SNe using the different features of its light curve and late-time spectrum. In this paper (Paper I), we investigate disk-driven outflows from the secondary BH, whereas possible signals from the primary BH will be discussed in Paper II.

Key words: accretion, accretion disks – binaries: close – black hole physics – gravitational waves – supernovae: general

1. Introduction

The advanced Laser Interferometer Gravitational-wave Observatory revealed the existence of black holes (BH) of ~30 M⊙ through the detection of gravitational waves (GWs) from mergers of binary black holes (BBHs; Abbott et al. 2016a, 2016b, 2016c). The origin of BBHs is under active debate, and several scenarios have been proposed, such as primordial black hole binaries (e.g., Nakamura et al. 1997; Mandic et al. 2016; Sasaki et al. 2016), multi-body interactions in star clusters (e.g., Sigurdsson & Hernquist 1993; Portegies Zwart & McMillan 2000; Rodríguez et al. 2016), and evolution of field binaries (e.g., Tutukov & Yungelson 1993; Kinugawa et al. 2014; Belczynski et al. 2016; Mandel & de Mink 2016; Marchant et al. 2016).

Future GW observations may provide the mass, spin, and redshift distributions of merging BBHs, which are useful to probe the environments, where BBHs are formed (e.g., Kushnir et al. 2016; Farr et al. 2017; Hotok ezaka & Piran 2017; Zaldarriaga et al. 2018). Searching for electromagnetic (EM) counterparts from merging BBHs is another way to study them. However, there is a substantial time gap, typically ~0.1–10 Gyr, between the merger events and the formation of BBHs, which makes it difficult to probe the environments of the BBH formation. Besides, simultaneous detections of EM counterparts and GWs are not guaranteed, because the possible EM signals considered so far require some specific conditions, such as the existence of a fossil disk (Murase et al. 2016; Perna et al. 2016; Ioka et al. 2017; Kimura et al. 2017a), or a BBH formation in anomalously dense environments, such as the inside of very massive stars (Loeb 2016; Dai et al. 2017) or active galactic nuclei (Bartos et al. 2017; Stone et al. 2017).

Instead, searching for the EM radiation from newborn BBHs would enable us more directly to probe the environment of BBH formation.

In this work, we suggest a new class of transient associated with the birth of a BBH system, which is a natural consequence of the evolution of the progenitor binary system from massive stars. The transients investigated in this work are not coincident with the GW emission at the BH–BH merger. Nevertheless, successful observations can provide important clues about the formation scenario of BBHs. We describe the basic binary evolution process and the possible outcomes in Section 2, where we consider two scenarios; one powered by the secondary BH and the other powered by the primary BH.

In this paper, we focus on the transient events driven by the secondary BH. We analytically estimate observable features of optical transients in Section 3 and the associated radio transients in Section 4. In Section 5, we summarize and discuss our results, including the observational prospects. The other type of transients powered by the primary BH are discussed in Kimura et al. (2017b, Paper II). We use the notation of A = A₁, 10^{4} throughout this work.

2. Consequences of the Evolutionary Scenario

According to the isolated binary evolution models (e.g., Belczynski et al. 2016), the heavier primary collapses to a BH earlier than the secondary, which forms a BH and main-sequence star binary. When the secondary evolves to a giant star, the binary separation decreases considerably during a common envelope evolution (Paczynski 1976; Webbink 1984).
After that, the secondary is expected to become a Wolf–Rayet star (WR), following the envelope ejection, and the BH–WR binary becomes a BBH after the gravitational collapse of the WR (e.g., Dominik et al. 2012). At the end of the binary evolution, after the massive secondary star has collapsed and the BBH has formed, its subsequent fate depends on the angular momentum of the secondary star. Although the angular momentum distribution of the secondary is highly uncertain, the spin of the secondary star may be tidally locked in a close binary system (Zahn 1977; Tassoul 1987). Kushnir et al. (2016) give the synchronization time as

\[
t_{\text{TL}} \sim 7.5 \times 10^4 \left( \frac{t_{\text{mer}}}{10^8 \text{year}} \right)^{17/8} \text{year},
\]

where \( t_{\text{mer}} = 5c^5a^4/(512G^3M_*^4) \sim 1.0 \times 10^8a_2^{-4}M_{1.5}^{-3} \text{year} \) is the GW inspiral time, \( M_* \) is the primary mass, and \( a \) is the binary separation. We assume the mass ratio \( q = 1 \) for simplicity and use \( M_* \sim 10^{1.5} M_\odot \) and \( a \sim 10^{12} \text{cm} \) for the purpose of an estimate, which indicates that \( t_{\text{TL}} \) is shorter than the typical lifetime of massive stars, \( t_{\text{life}} \sim 10^6 \text{yr} \). However, for a low mass \( M_* \sim 10 M_\odot \) or large separation \( a \sim 3 \times 10^{12} \text{cm} \), \( t_{\text{TL}} > t_{\text{life}} \) is possible. We caution that this timescale has significant uncertainties caused by the strong dependence on the detailed stellar structure, especially the size of the convective region (Kushnir et al. 2016). The size of the star also affects this timescale.

When \( t_{\text{TL}} < t_{\text{life}} \), the spin period of the secondary is synchronized to its orbital period. The spin angular momentum of the WR is high enough to prevent the WR from directly collapsing to a BH. The outer region of the WR forms an accretion disk around the newborn secondary BH. The accretion rate is high enough to produce radiation-driven powerful outflows (e.g., Ohsuga et al. 2005; Jiang et al. 2014), leading to a tidally locked secondary supernova (TLSSN, see Figure 1 for the schematic picture). The kinetic energy of this outflow is so large that we can expect a radio afterglow.

In the opposite case when \( t_{\text{TL}} > t_{\text{life}} \), the WR is likely to spin slowly enough to collapse to a BH directly. When the WR collapses, the outer envelope of the WR is ejected due to energy losses by neutrinos (Nadezhin 1980). The primary BH accretes the ejected material, and may produce powerful outflows owing to its high accretion rate. This outflow energizes the ejecta and could lead to a primary-induced accretion transient (PIATs), which is discussed in the accompanying paper (Paper II).

A disk-driven outflow can produce a super-luminous supernova (Dexter & Kasen 2013), a hypernova (MacFadyen & Woosley 1999), and an optical transient during a single BH formation (Kashiyama & Quataert 2015) and BH mergers (Murase et al. 2016). TLSSNe and PIATs provide different examples associated with a newborn BBH.

3. Optical Emission from TLSSNe

We consider a tidally synchronized binary system where the spin period of the secondary is synchronized to the orbital period of the primary. As binary parameters, we choose the mass of the WR, \( M_* \sim 10^{1.5} M_\odot \), the radius of the WR, \( R_\text{e} \sim 10^{11} \text{cm} \), the binary separation, \( a \sim 10^{12} \text{cm} \), and the mass ratio, \( q = M_/M_\text{BH} = 1 \). This parameter set satisfies \( t_{\text{TL}} < t_{\text{life}} \) and is consistent with stellar evolution models (Schaerer & Maeder 1992). The spin angular velocity is synchronized to the orbital motion, which is estimated to be \( \omega_i = \sqrt{2GM_*/a^3} \sim 9.2 \times 10^{-4}M_{1.5}^{-1/2}a_2^{-3/2} \text{s}^{-1} \). This value is so high that the outer part of the stellar material cannot fall toward the BH directly. Thus, an accretion disk is formed.

\[ \text{Reference:} \]  
\[ \text{Footnote:} \]

5 The centrifugal and Colioris forces do not affect a disk formation process when the WR star collapses even when the binary separation is considerably close.

6 There is some uncertainty for radii of WR stars. A relation \( R_\text{e} \sim 7 \times 10^{10} (M_/10 M_\odot)^{0.7} \text{cm} \) is obtained by stellar evolution models (Schaerer & Maeder 1992; Kushnir et al. 2016), while \( R_\text{e} \sim 2 \times 10^{11} \text{cm} \) are proposed from an atmospheric model (Crowther 2007). Besides, the radius is larger for lighter secondaries, \( R_\text{e} \sim 10^9 \text{cm} \) for \( M_* \sim 5 M_\odot \), according to a binary evolution model (Yoon et al. 2010).

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Figure 1. Schematic picture of tidally locked secondary supernovae (TLSSNe). (i) A WR is synchronized via the tidal force before its collapse. The inner part forms a secondary BH, while the outer material forms a disk around the BH. (ii) An ejecta is launched by the disk-driven outflow. (iii) Thermal photons diffuse out from the ejecta.
around the newborn BH. This produces a massive outflow, which leads to a TLSSN.

When the secondary collapses, the outer material of the secondary at a cylindrical radius \(r\) is at the centrifugal radius, 
\[ r_c(R) = \frac{\omega^2}{G M_{\ast}} \approx 2 \omega^2 R^3, \]
where \(M_{\ast} \sim M_{\ast, 15} \) is the mass enclosed inside the spherical radius \(R = 3M_{\ast} \).

Setting \(r_c = 6GM_{\ast}/c^2\), we obtain the critical radius for disk formation:
\[ R_c \approx \left(\frac{3GM_{\ast}}{c^2}\right)^{1/4} \approx 6.1 \times 10^{30} M_{\ast, 15}^{1/4} \text{ cm}. \]  

(2)

Since \(R_c > R_{\ast}\) for our parameter choice, the outer material can form a disk. The density profile of WRs can be expressed as a polytropic sphere of index \(\sim 2.5\)–3.5 (Kushnir et al. 2016). The outer region \((R \gg R_{\ast}/2)\) of a polytropic n can be fitted as \(\rho_{\text{env}} \approx \rho_n(R_{\ast}/R)^{-n}\) (Matzner & McKee 1999), where \(\rho_n \approx A_n M_{\ast}/(4\pi R_{\ast}^2)\) and \(A_n \approx 3.9\) is a numerical constant that depends on the polytropic index (we use \(n = 3\)). This expression can reproduce the polytrope within errors of a few percent, except for the very outer edge, which does not affect the result.

The disk mass is then estimated to be
\[ M_d = 4\pi \int_{R_{\ast}}^{R_c} r^2 d \omega \int_0^{\sqrt{R_c^2 - r^2}} dz \rho_n \left(\frac{R_c^2}{r^2 + z^2} - 1\right)^{n/2} \approx A_n I_d M_{\ast} \approx 0.41 M_{\ast, 15} L_{\ast, 5} \text{ M}_{\odot}. \]  

(3)

where \(I_d = \int_0^1 dx \int_0^{(1-x)^{1/2}} dy (1 - y^2)^{n/2} \approx 3.3 \times 10^{-3}\) and \(x_{\text{cr}} = R_{\text{cr}}/R_{\ast}\). Note that \(I_d\) is a strong function of \(x_{\text{cr}}\) that depends on \(M_{\ast}\), so that the dependence of \(M_d\) on \(M_{\ast}\) is not simple. The outer material falls to the disk in a free-fall time, 
\[ t_{\text{ff}} \approx \left(3M_{\ast}/(GM_{\ast})\right)^{1/4} \approx \frac{7.5}{4} \times 10^{31} \text{ M}_{\ast, 15}^{1/2} \text{ R}_{\ast, 11}^{1/2} \text{ s}. \]

Then, the mass accretion rate is estimated to be
\[ \dot{M}_{\text{env}} \approx \frac{M_d}{t_{\text{ff}}} \approx 7.5 \times 10^{-4} M_{\ast, 15}^{1/2} \text{ R}_{\ast, 11}^{3/2} M_{\ast, 0.39} \text{ M}_{\odot}^{-1}\text{s}^{-1}. \]  

(4)

This accretion rate is much higher than the Eddington rate, 
\[ \dot{M}_{\text{Edd}} = L_{\text{Edd}}/c^2 \approx 2.2 \times 10^{-15} M_{\ast, 15} \text{ M}_{\odot}^{-1}s^{-1}, \]
while it is much lower than the critical mass accretion rate for neutrino cooling, 
\[ \dot{M} \sim 1 \text{ M}_{\odot}^{-1}s^{-1}. \]

(Popham et al. 1999; Kohri & Mineshige 2002). Then, the physical state of the accretion flow is expected to be the advection dominant regime, where the outflow is likely to be produced (Narayan & Yi 1994; Blandford & Begelman 1999; Kohri et al. 2005). The wide-angle outflow production is also commonly seen in numerical simulations for the central engine of gamma-ray bursts (GRBs) (MacFadyen & Woosley 1999; Fernández & Metzger 2013). The outflow luminosity is estimated to be
\[ L_w \approx \frac{1}{2} \eta_w \dot{M}_{\text{env}} V_w^2 \approx 2.4 \times 10^{49} \text{ M}_{\ast, 15}^{1/2} \text{ R}_{\ast, 11}^{3/2} \dot{M}_{\ast, 0.39} \eta_{w, 0.39} \text{ V}_{10}^2 \text{ erg s}^{-1}, \]  

where we assume that \(\eta_w \sim 1/3\) of the accreting material is ejected as the outflow with velocity \(V_w \sim 10^8 \text{ cm s}^{-1}\). Although these parameters related to the outflow are highly uncertain, these values of \(\eta_w\) and \(V_w\) are consistent with the recent simulation and observation results (Hagino et al. 2015; Takahashi et al. 2016; Narayan et al. 2017).\(^7\)

The duration of the outflow is comparable to the free-fall time, since the accretion time after the disk formation is much shorter than the free-fall time. The total mass of the outflow is \(M_w = \eta_w M_d \approx 0.14 M_{\ast, 0.39} \text{ M}_{\odot}\) and the total energy is
\[ E_w \approx \frac{1}{2} \eta_w \dot{M}_d V_w^2 \approx 1.3 \times 10^{52} M_{\ast, 0.39} \text{ erg}. \]  

(5)

We assume that the outflow occurs at the escape velocity, and is launched at \(v_{\text{p}} \approx 2GM_{\ast}/V_w \approx 8.4 \times 10^7 \text{ M}_{\ast, 15}^{-1/2} \text{ cm} \). Assuming that the radiation energy is comparable to the kinetic energy at the launching point, the temperature at that point is
\[ T_{\text{p}} \approx \left(\frac{M_w V_w}{8\pi \sigma r_{\text{p}}^2}\right)^{1/4} \approx \left(\frac{1.4 \times 10^{54} \text{ M}_{\ast, 15}^{3/8} \text{ R}_{\ast, 11}^{3/8} \dot{M}_{\ast, 0.39}^{1/4} \text{ V}_{10}^{1/4}}{\text{K}}\right). \]

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(7)

The outflow is expected to be
\[ E_{\text{out}} = \int_{r_{\text{p}}}^{R_d} d\tau \frac{v_{\text{p}}}{v_{\text{in}}} \text{ erg} \approx 3 \eta_w \dot{M}_d \frac{GM_{\ast}}{R_{\ast}}. \]  

(8)

The bulk of photons are trapped inside the outflow, but a small fraction of the photons can escape from the outflow through a diffusion process. The luminosity of the escaping photons is estimated to be
\[ L_{\text{ph,0}} \approx L_{\text{in,0}} \approx 9.0 \times 10^{41} \text{ M}_{\ast, 15}^{1/2} \text{ R}_{\ast, 11}^{2/3} V_{10} \text{ ergs s}^{-1}. \]  

(9)

where \(\tau_{\text{ph,0}} \approx 9 \eta_n \dot{M}_d/(4\pi^3 R_{\ast, 0.0} c)^{1/2}\) is the photon diffusion time at \(t = t_{\text{ff}}\). Since the optical depth is large enough, the escaping photons have a Planck spectrum with an effective temperature of
\[ T_{\text{eff,0}} \approx \frac{L_{\text{ph,0}}}{4\pi \sigma T_{\text{ph,0}}^4} \approx 8.5 \times 10^4 \text{ M}_{\ast, 15}^{3/8} \text{ R}_{\ast, 11}^{5/8} \text{ V}_{10}^{1/4} \text{ K}. \]  

(10)

Note that these simulations and observations are for the cases with \(M \sim 10^{-2} M_{\text{Edd}}\). The values of \(\eta_w\) and \(V_w\) for \(M \sim 10^{39} M_{\text{Edd}}\) should be investigated in the future.
estimated to be (Popov 1993)

\[ t_\text{i} \approx \left( \frac{L_{\text{ph},0}}{4\pi\sigma T_{\text{ion}}^4 V_w} \right) \simeq 3.6 \times 10^4 M_{*,15}^{1/4} R_{*,11}^{1/4} V^{-3/2}_{10} \text{ s}. \]  

(11)

The position of the recombination surface is given by (Popov 1993)

\[ R_\text{i} \approx \frac{V_w^2}{\pi} \left[ t_\text{i} \left( 1 + \frac{t_\text{i}^2}{3a_w^2} \right) - \frac{t^4}{3a_w^2} \right], \]

(12)

where \( t_\text{i} = \sqrt{2} t_{\text{ph},0} R_{w,0}/V_w \) is the photon breakout time without the recombination surface (Arnett 1980). Here, we assume \( t_\text{i} \ll t_\text{w} \) when estimating \( t_\text{i} \). Since the opacity of the neutral gas is very small (e.g., Kleiser & Kasen 2014), the photosphere is equal to the recombination surface. Thus, the effective temperature is \( T_{\text{eff}} = T_{\text{ion}} \), and the luminosity is estimated to be \( L_{\text{ph}} = 4\pi\sigma T_{\text{ion}}^4 R_\text{i}^2 \). Since \( t_\text{i}^2 \ll a_w^2 \) is satisfied for the parameter range of our interest, the luminosity has a maximum value of

\[
L_{\text{pk}} \approx 4\pi\sigma T_{\text{ion}}^4 V_w^2 \left( \frac{3t_{\text{i}}^2 a_w^2}{4} \right) \simeq 5.4 \times 10^{42} M_{*,15}^{1/6} R_{*,11}^{1/6} M_{\text{d},0,39}^{2/3} V_{10}^{11/3} \text{ erg s}^{-1},
\]

(13)

and the time of the peak luminosity is

\[
t_{\text{pk}} \approx \left( \frac{3t_{\text{i}}^2 a_w^2}{4} \right)^{1/3} \simeq 8.7 \times 10^4 M_{*,15}^{1/12} R_{*,11}^{1/12} M_{\text{d},0,39}^{2/3} V_{10}^{-5/6} \text{ s}. \]

(14)

For \( t > t_{\text{pk}} \), \( R_\text{i} \) rapidly decreases, and accordingly the luminosity decays quickly. The optical depth for the outflow becomes lower than unity in this phase. The approximate use of the Planck distribution would become inaccurate, especially at \( \tau \lesssim 1-10 \). To study the features of the decay phase, a more careful treatment of thermalization processes would be required.

The evolution of \( T_w \) and \( L_{\text{ph}} \) are shown in the upper panel of Figure 2, where we use the fiducial parameter set \( (M_*=30 M_\odot, R_* = 10^{11} \text{ cm}, a = 10^{12} \text{ cm}, \eta_w = 10^{-0.5}, V_w = 10^{10} \text{ cm s}^{-1}) \). The lower panel of the figure shows the evolution of the absolute AB magnitude in the \( U (365 \text{ nm}), V (550.5 \text{ nm}) \), and \( R (658.8 \text{ nm}) \) bands. Since the outflow parameter is uncertain, we show the results for the case with \( \eta_w = 10^{-1.5} \) and \( V_w = 10^{9.5} \text{ cm s}^{-1} \) in Figure 3 for comparison. We can see that the optical band light curves rapidly become bright in several hours, remain bright and slowly varying for days, and then rapidly fade on a timescale of hours. The peak magnitude range is \(-14 \sim -17\), which is similar to that of usual type II or type Ib/Ic SNe. However, their shorter durations are useful for distinguishing TLSSNe from the usual SNe. Spectroscopic observations can also discriminate TLSSNe from macronovae/kilonovae, since TLSSNe will show strong helium lines while macronovae/kilonovae are not expected to show such lines. We note that these TLSSNe are bright and short duration, compared to the PIAT events discussed in Paper II.

Interestingly, \( t_{\text{pk}} \) and \( L_{\text{pk}} \) do not have a strong dependence on the parameters. However, the occurrence of TLSSNe is sensitive to the value of \( x_{\text{cr}} = R_{\text{cr}}/R_\odot \simeq 0.61 M_{*,15}^{1/4} R_{*,11}^{-1/3} a_{\text{d},0,39}^{2/3} \). For \( x_{\text{cr}} > 1 \), an accretion disk is not formed, which leads to a result similar to the PIATs that we discuss in Paper II. For \( x_{\text{cr}} \ll 1 \), a significant fraction of the stellar material falls onto the disk, and the newborn BH has a large spin, probably resulting in a GRB. See Section 5 for a discussion about the possible relation between GRBs and TLSSNe.

4. Radio Afterglows of TLSSNe

Outflow-driven transients may lead to radio afterglows (see Kasliyama et al. 2017 for details; see also Murase et al. 2016). We briefly discuss this possibility here (see Chevalier 1998; Nakar & Piran 2011, for supernovae and neutron star mergers). The deceleration radius and time is estimated to be

\[
R_{\text{dec}} \approx \left( \frac{3M_w}{4\pi n_{\text{ext}} m_p} \right)^{1/3} \simeq 5.0 \times 10^{18} M_{\text{d},0,39}^{1/3} V_{10}^{-1/3} n_{-1}^{-1/3} \text{ cm},
\]

(15)

\[
t_{\text{dec}} \approx \frac{R_{\text{dec}}}{V_w} \simeq 5.0 \times 10^8 M_{\text{d},0,39}^{1/3} V_{10}^{-1/3} n_{-1}^{-1/3} \text{ s},
\]

(16)
where \( n_{\text{ext}} = 0.1 \, \text{cm}^{-3} \) is the number density of the circumbinary medium. The deceleration time can be shorter for smaller \( \eta_\nu \) and larger \( V_\nu \). After the deceleration time, the evolution of the outflow is represented by the self-similar solution,

\[
R = R_{\text{dec}}(t/t_{\text{dec}})^{2/5}(t \geq t_{\text{dec}}),
\]

\[
V = 0.4 V_\nu (t/t_{\text{dec}})^{-3/5}(t \geq t_{\text{dec}}).
\]

We estimate the physical quantities around \( t \sim t_{\text{dec}} \) using \( V \sim V_\nu \) and \( R \sim R_{\text{dec}} \). The magnetic field is estimated to be

\[
B = (9 \pi m_e n_{\text{ext}} \epsilon_B \nu^2) \approx 1.4 V_{\nu0} n_{-1}^{1/2} \epsilon_{-2}^{1/2} \, \text{mG},
\]

where \( \epsilon_B \) is the energy fraction of the magnetic field. The minimum Lorentz factor of electrons is approximately

\[
\gamma_m \approx \frac{\zeta_e}{2} \frac{m_p}{m_e} \left( \frac{V}{c} \right)^2 \approx 41 V_{\nu0}^{2} \zeta_{-0.4},
\]

where \( \zeta_e \approx (p - 2) \epsilon_e / (p - 1) f_e \) \( \approx 0.4 \), \( \epsilon_e \) is the energy fraction of the nonthermal electron, \( f_e \approx 0.1 \) is the number fraction of nonthermal electron, and \( p \) is the spectral index of the nonthermal electrons. The cooling Lorentz factor is \( \gamma_c \approx 6 \pi m_e c / (\sigma_T B^2) \approx 7.4 \times 10^5 M_{d,-0.39}^{-1/3} \eta_{-0.5}^{-1/5} V_{\nu0}^{-1} n_{-1}^{-2/3} \epsilon_{-2}^{-1} \), where \( \sigma_T \) is the Thomson cross section. Since \( \gamma_m \ll \gamma_c \), the synchrotron spectrum is in the slow cooling regime. The synchrotron frequencies for the electrons of \( \gamma_m \) and \( \gamma_c \) are

\[
\nu_m \approx \frac{\gamma_m^2 e B}{2 \pi m_e c} \approx 6.7 \times 10^6 V_{\nu0}^{-1} n_{-1}^{1/2} \epsilon_{-2}^{1/2} \zeta_{-0.4},
\]

\[

\nu_c \approx 2.2 \times 10^{15} M_{d,-0.39}^{-2/3} \eta_{-0.5}^{-1/5} V_{10}^{-3} n_{-1}^{-5/6} \epsilon_{-2}^{-3/2} \text{Hz}.
\]

If we ignore synchrotron self-absorption (SSA), the synchrotron spectrum has a peak at \( \nu_m \) and its flux is

\[
F_{\nu,m} \approx \frac{P_{\nu0} \eta_\nu}{4 \pi v_\nu d_L^2} \approx 22 M_{d,-0.39}^{-1/2} V_{\nu0}^{1/2} n_{-1}^{1/2} f_{-1} d_L^{-2} \text{mJy},
\]

where \( P_{\nu0} \approx \gamma_m^2 \sigma_T c B^2 / (6 \pi \nu) \) is the synchrotron radiation power per electron, \( \eta_\nu \approx 4 \pi R^2 n_{\text{ext}} f_e / 3 \) is the total nonthermal electron number, and \( d_L \approx 10^{27} \text{cm} \) is the luminosity distance. Since \( F_{\nu0} \propto \nu^{(1-\nu)/2} \) for \( \nu_m < \nu < \nu_c \) without SSA, the observed flux at frequency \( \nu_{\text{obs}} \) is estimated to be

\[
F_{\nu,\text{obs}} \approx F_{\nu,m} \left( \frac{\nu_{\text{obs}}}{\nu_m} \right)^{(1-\nu)/2} \approx 0.15 \nu_{\text{obs}}^{-2/3} M_{d,-0.39}
\times \eta_{-0.5} V_{10}^{5} n_{-1}^{4} f_{-1} d_L^{-2} \text{mJy},
\]

where we use \( p = 3 \) to estimate the value. Since the sensitivity of current radio surveys is around 0.1 mJy, it is possible to detect this radio emission. If the optical transient discussed in Section 3 is observed, the deeper radio follow-up observation with a sensitivity of around \( \mu Jy \) can be performed. In this case, we can expect detection of the radio signal even with much lower \( \eta_\nu \). The deceleration time is shorter for lower \( \eta_\nu \), which helps the coincident detection. Using Equations (17) and (18), we obtain \( F_{\nu,\text{obs}} \propto t^2 \) for \( t < t_{\text{dec}} \) and \( F_{\nu,\text{obs}} \propto t^{(21/15)\nu/10} \) for \( t > t_{\text{dec}} \). Note that \( F_{\nu,\text{obs}} \) has a strong dependence on \( V_\nu \), \( F_{\nu,\text{obs}} \propto V_\nu^4 \) for \( p = 3 \). Thus, just a few times lower \( V_\nu \) would make it difficult to detect the radio afterglow.

The optical depth for SSA is estimated to be

\[
\tau_\nu \approx A_p \frac{\sigma_T n_{\text{ext}} R^2 (\nu/\nu_m)^{(p+4)/2}}{(B^4 m^5)},
\]

where \( A_p \) is a function of \( p \) (\( A_p = 26.31 \) for \( p = 3 \), see Murase et al. 2014). The SSA frequency at which \( \tau_\nu = 1 \) is estimated as

\[
\nu_\nu \approx \frac{A_p \sigma_T n_{\text{ext}} R^2}{B^4 m^5} 2^{(p+4)} \nu_m \approx 7.1 \times 10^5 M_{d,-0.39}^{2/3} V_{10}^{5} n_{-1}^{4} f_{-1} d_L^{-2} \epsilon_{-2}^{-1} \text{Hz}.
\]

This frequency evolves as \( \nu_\nu \propto t^{2(3p+2)/(p+4)} \) for \( t < t_{\text{dec}} \) and \( \nu_\nu \propto t^{-4(3p+2)/(p+4)} \) for \( t > t_{\text{dec}} \) (Nakar & Piran 2011). If we focus on \( \nu_{\text{obs}} > \nu_\nu \), we can ignore the effect of SSA. Since we typically expect \( \nu_\nu > \nu_{\text{obs}} \), the spectrum is modified by SSA as \( \nu_\nu \propto \nu^{2p+2} \eta_{-0.5}^{-1/5} V_{10}^{-1} n_{-1}^{-5/6} \epsilon_{-2}^{-1} \text{Hz} \).

### 5. Summary and Discussion

We investigated outflow-driven transients from newborn BBHs formed from BH–WR binaries, within the context of isolated binary evolution scenarios. When the binary separation is small or the binaries are massive enough, the spin period of the WR is synchronized to the orbital period. When the WR collapses to a BH, the outer region of the WR has such a high angular momentum that an accretion disk is formed around a
newborn secondary BH. This results in an energetic outflow of kinetic energy of $\sim 10^{52} \text{ erg}$ for $\eta_0 \sim 10^{-3}$, leading to a TLSSN whose bolometric luminosity can be $\sim 10^{42} - 10^{43} \text{ erg s}^{-1}$. Its optical band absolute magnitude reaches approximately $-17$, with a duration of around a day. Transient radio emission can also be expected, owing to the large amount of kinetic energy involved.

When the binary separation is larger or the stellar mass is lower, the tidal synchronization may not occur and the spin of the secondary is likely to slow down. Even in this case, a fraction of the outer material of the secondary is ejected when the secondary collapses to a BH. This ejected material is expected to be accreted by the primary BH, and a powerful outflow is produced, leading to a PLAT. We discuss this type of transient in the accompanying paper (Paper II).

The TLSSNe can be distinguished from usual SNe by their shorter duration, and from macronovae/kilonovae by their strong helium lines. The light curves of TLSSNe are consistent with some of the rapid transients observed (Drout et al. 2014; Tanaka et al. 2016) on the basis of their timescale (around a day) and absolute magnitude (approximately $-16$, although other phenomena, such as shock breakouts from cooling envelopes (e.g., Waxman et al. 2007) or the outflow-driven transient from single BH formation (Kashiyama & Quataert 2015) could also appear similar.

The current optical surveys with a sensitivity of $\sim 21$ mag, such as Pan-STARRS (Hodapp et al. 2004), PTF (Law et al. 2009), and KISS (Morokuma et al. 2014), imply a detectability distance for TLSSNe of $\sim 200$ Mpc. Assuming that the event rate of TLSSNe is similar to the merger rate of BBHs, $10^{3} - 10^{4} \text{ Gpc}^{-3} \text{ yr}^{-1}$ (Abbott et al. 2016a, 2017), the event rate within the sensitivity range is $0.3 - 7 \text{ yr}^{-1}$. Thus, the current surveys could detect this type of transient in the near future. However, we should note that the event rate of TLSSNe has substantial uncertainties, related to the binary evolution and the outflow from a super-Eddington accretion flow. Future projects, e.g., the Large Synoptic Survey Telescope (LSST; Abell et al. 2009), would be able to detect them or put a meaningful limit on the event rate.

The stellar wind, the ejected stellar envelopes during the common envelope phase, and/or the supernova impostors can significantly pollute the circumbinary medium (e.g., Smith et al. 2011; Beclczynski et al. 2016). This circumbinary matter can be more massive than the outflow of TLSSNe. Thus, they could affect both the optical and radio light curves of TLSSNe.

Since the core of the WR can be radiative, it is possible that the core rotates faster than the envelope. If the core rotates sufficiently fast, it forms a BH with an accretion disk $\sim 1$ second after the core-collapse (O’Connor & Ott 2011). The mass accretion rate is large enough to synthesize some amount of radioactive nuclei around this secondary BH (Pruet et al. 2003; Fujimoto et al. 2004), which may produce another type of supernova/hypernova-like transient powered by the radioactive decay of nuclei. Even relativistic jets that could lead to a GRB may be launched, and during the jet propagation phase, the radioactive nuclei can be synthesized in the shocked envelope, leading to another energy source for supernova-like emission (Barnes et al. 2017). Different energy sources may coexist in the tidally locked system we here consider, so that optical emission from TLSSNe may be powered by either the thermal emission from the disk outflow itself or regenerated emission from radioactive nuclei.

If the core commonly rotates very fast, such a binary system may be responsible for long GRBs (see Fryer & Heger 2005). Indeed, the true event rate of GRBs after beaming correction is consistent with the expected event rate of TLSSNe. In addition, the tidally locked system can naturally produce late-time central engine activity for plateau emission or X-ray flares in GRB afterglows, because the outer envelope accretes onto the BH $\sim 10^{3} \text{ s}$ after the core accretion that is attributed to the GRB prompt emission. Note that a disk state for a typical TLSSN discussed in this paper is different from that of a collapsar disk discussed in the context of GRBs (MacFadyen & Woosley 1999). The disk in a TLSSN has much lower temperature than that of a collapsar disk, so a TLSSN disk cannot produce a jet through the neutrino annihilation (Eichler et al. 1989; Popham et al. 1999). If the secondary BH has a high spin and global magnetic field, the magnetic jet can be produced (Blandford & Znajek 1977; Komissarov 2004; Toma & Takahara 2016). Although the jet power seems too low to produce a typical long GRB, it may be observed as an ultra-long GRB if the jet is directed at the Earth (Quataert & Kasen 2012; Woosley & Heger 2012). Since a wide-angle outflow can simultaneously be produced (MacFadyen & Woosley 1999), we may also observe a TLSSN.

An accretion disk around a BH in a BBH is left over after the transients considered here. A few years later, this disk is expected to become a fossil disk, in which the angular momentum transport is inefficient, due to radiative cooling (e.g., Perna et al. 2014, 2016; Kimura et al. 2017a). Since such fossil disks can remain for millions of years, a possible outcome from them would be electromagnetic counterparts of the GWs from the eventual BH mergers (Murase et al. 2016; de Mink & King 2017a; Kimura et al. 2017a).

Besides the transients discussed here, which involve a WR companion, there are likely to be other formation channels of BBHs through binary evolution, where the progenitor consists of a BH and a blue-super giants (BSGs) or red super giants (RSGs). Since BSGs and RSGs have larger radii than WRs, $x_r < 1$ is easily satisfied. In this sense, TLSSNe are likely in BH–BSG and BH–RSG binaries. However, whether the spin is tidally synchronized or not depends strongly on the internal structure of the secondary (Kushnir et al. 2016). Also, the tidal force from the primary distorts the WR star to nonspherical shape, which could affect the stellar structure. A more accurate modeling will require solving the stellar evolution in detail. Due to these uncertainties related to the stellar structure as well as the outflow properties resulting from the super-Eddington accretion, it is currently difficult to derive a meaningful luminosity distribution for such transients.

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**References**

Abbott, B. P., Abbott, R., Abbott, T. D., et al. 2016a, PhRvX, 6, 041015
Abbott, B. P., Abbott, R., Abbott, T. D., et al. 2016b, PhRvL, 116, 241103
Abbott, B. P., Abbott, R., Abbott, T. D., et al. 2016c, PhRvL, 116, 061102
Abbott, B. P., Abbott, R., Abbott, T. D., et al. 2016d, PhRvL, 116, 101104
