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Charged Compact Binary Coalescence Signal and Electromagnetic Counterpart of Plunging Black Hole–Neutron Star Mergers

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Abstract

If at least one of the members of a compact binary coalescence is charged, the inspiral of the two members would generate a Poynting flux with an increasing power, giving rise to a brief electromagnetic counterpart temporally associated with the chirp signal of the merger (with possibly a small temporal offset), which we term as the charged compact binary coalescence (cCBC) signal. We develop a general theory of cCBC for any mass and amount of charge for each member. Neutron stars (NSs), as spinning magnets, are guaranteed to be charged, so the cCBC signal should accompany all NS mergers. The cCBC signal is clean in a black hole (BH)–NS merger with a small mass ratio ($q \equiv m_2/m_1 < 0.2$), in which the NS plunges into the BH as a whole, and its luminosity/energy can reach that of a fast radio burst if the NS is Crab-like. The strength of the cCBC signal in Extreme Mass Ratio Inspiral Systems is also estimated.

Key words: gravitational waves – radiation mechanisms: general – stars: black holes – stars: neutron

1. Introduction

The discovery of the neutron star–neutron star (NS–NS) merger gravitational wave (GW) event GW170817 (Abbott et al. 2017b) and its associated electromagnetic (EM) signals, including the short gamma-ray burst (GRB) 170817A (Abbott et al. 2017a; Goldstein et al. 2017; Zhang et al. 2018), the "kilonova" AT2017gfo (Coulter et al. 2017; Villar et al. 2017), as well as the broadband (from radio to X-rays) nonthermal "afterglow" (Hallinan et al. 2017; Troja et al. 2017; Lyman et al. 2018; Margutti et al. 2018; Piro et al. 2019), formally ushered in the era of the GW-led "multi-messenger" astrophysics.

Neutron star mergers (including NS–NS and black hole (BH)–NS mergers) are widely believed to be bright EM emitters. This is because such systems have plenty of neutron-rich matter outside the horizon of the BH (if formed) in the merger remnant. Gravitational energy and/or nuclear energy are released as the matter is accreted into the BH to produce a short GRB (Eichler et al. 1989; Mészáros & Rees 1992; Narayan et al. 1992), or as the outgoing ejecta undergo the r-process nucleosynthesis and the subsequent β-decay to power a kilonova (Li & Paczyński 1998; Metzger et al. 2010; Barnes & Kasen 2013; Kasen et al. 2017). If the post-merger remnant is a long-lived neutron star rather than a BH, additional energy release is possible with the expense of the spin energy and magnetic energy of the neutron star and its pulsar wind (Zhang 2013; Gao et al. 2013; Yu et al. 2013; Metzger & Piro 2014; Piro et al. 2019).

For binary BH–BH mergers or "plunging" BH–NS mergers with a mass ratio $q < 0.2$ (Shibata et al. 2009), it is widely believed that no bright EM emission is expected from the systems, because matter is contained within the horizon of the post-merger BH. This was why the claimed γ-ray counterpart, GW150914-GBM, associated with the first BH–BH merger event was regarded as being controversial (Connaughton et al. 2016, 2018; Greiner et al. 2016), and why most theoretical models proposed to interpret the event need to introduce contrived physical conditions to maintain enough matter outside the horizon right after the merger (e.g., Loeb 2016; Perna et al. 2016, see Dai et al. 2017; Kimura et al. 2017).

Zhang (2016) suggested that instead of maintaining matter outside the horizon, one may maintain an EM field in at least one member of the compact binary coalescence (CBC) in order to power an EM counterpart of the CBC. He found that the Poynting-flux luminosity from the system sharply rises right before the merger, which hereafter we refer to the charged compact binary coalescence (cCBC) signal. He hypothesized that at least one BH of the GW150914 system may be significantly charged. He further suggested that with a smaller charge, BH–BH merger systems may account for a fraction of non-repeating fast radio bursts (FRBs), the mysterious milli-second radio bursts at cosmological distances (Lorimer et al. 2007; Thornton et al. 2013). The charged BH–BH mergers were invoked in several other studies to account for the EM counterparts of BH–BH mergers or FRBs (Liebling & Palenzuela 2016; Liu et al. 2016; Deng et al. 2018; Fraschetti 2018; Levin et al. 2018). Whether BHs can sustain a significant amount of charge is still an open question. Zhang (2016) argued that a spinning charged (Kerr–Newmann) BH can in principle carry a force-free magnetosphere. Like magnetized spinning NSs (pulsars), they can sustain a global charge due to the spatial distribution of the charge density demanded to maintain a co-rotating force-free magnetosphere.

In this Letter, we discuss the cCBC signal in general. Because NSs are spinning magnets, they are guaranteed to be charged (Michel 1982; Pétri 2012). Even if BHs may not maintain a large enough charge long enough until the merger occurs, NS–NS and BH–NS merger systems should have at least one member (the NS) charged, and the physical process delineated in Zhang (2016) should apply. Given the relatively small charge sustained by NSs, the cCBC signal is insignificant in NS–NS merger systems and most BH–NS systems with a relatively large mass ratio $q \equiv m_2/m_1$, where m_1 and $m_2 < m_1$ are the masses of the two members in the CBC. This is because the accretion-powered short GRB signal in these systems is orders of magnitude brighter than the cCBC signal. Furthermore, the dynamical ejecta launched during the merger places a significant opacity over a very large solid angle. Certain cCBC signals, e.g., an FRB, would be subject to absorption and may not escape in most solid angles. On the other hand, for BH–NS merger systems with $q < 0.2$ (e.g., Shibata et al. 2009), the neutron star does not undergo tidal disruption before the merger, which plunges into the BH entirely. Previously, it has been believed that such systems may not produce EM signals (e.g., Bartos et al. 2013). Here we suggest that these are ideal systems to observe the cCBC signal. A general theory of cCBC is presented in Section 2. The strength of the signal for plunging BH–NS systems is discussed in Section 3. The case of extreme mass ratio inspiral (EMRI) systems is discussed in 4. The results are summarized in Section 5 with some discussion.

2. General Theory of Charged CBC

Most generally, we consider two members in the CBC that are characterized by $(m_1, \hat{q}_1)$ and $(m_2, \hat{q}_2)$, respectively, where m_i is the mass of the member $i = 1, 2$, and

$$\hat{q}_i \equiv Q_i/Q_{c,i}$$

is the absolute charge Q_i divided by its critical charge defined by (Zhang 2016)

$$Q_{c,i} \equiv 2\sqrt{G}m_i,$$

and G is the gravitational constant. Besides the total mass $M \equiv m_1 + m_2$ and the mass ratio $q = m_2/m_1$, one can define three additional masses from m_1 and m_2:

$$\text{Reduced mass:}\quad M_r = \frac{m_1 m_2}{m_1 + m_2}; \tag{1}$$

$$\text{Chirp mass:}\quad M_c = M_r^{3/5} M^{2/5}; \tag{2}$$

$$\text{Horizon mass:}\quad M_h = M_r^{2/5} M^{3/5}. \tag{3}$$

The first two masses are commonly used in the GW community, and the meaning of the third mass (M_h) will be introduced shortly.

2.1. GW Luminosity

For easy comparison with the EM luminosities presented in later subsections, it is informative to write down the GW luminosity in several different forms. The first two forms are the standard expressions (e.g., Maggiore 2008)

$$L_{\rm GW} = \frac{32}{5}\frac{G^4}{c^5}\frac{\mu^2 M^3}{a^5} f(e), \tag{4}$$

$$= \frac{32}{5}\frac{c^5}{G}\left(\frac{G M_c \omega_s}{c^3}\right)^{10/3} f(e), \tag{5}$$

where c is speed of light, a is the semimajor axis, $\omega_s \equiv (GM/a^3)^{1/2}$ is the orbital angular frequency, and

$$f(e) = \frac{1 + \frac{73}{24}e^2 + \frac{37}{96}e^4}{(1 - e^2)^{7/2}} \tag{6}$$

is a correction factor introduced by the orbital eccentricity e, which equals unity when $e = 0$ (circular orbit). Notice that

$$\frac{c^5}{G} \simeq 3.63 \times 10^{59}\ \text{erg s}^{-1} \tag{7}$$

is a luminosity unit defined by fundamental constants.

As ω_c is directly related to the frequency of GWs, $\omega_{\rm gw} = 2\omega_s$, the chirp mass M_c can be directly connected to the observed quantities according to Equation (5). For the purpose of our study, the GW luminosity can be expressed in another form:

$$\begin{aligned} L_{\rm GW} &= \frac{1}{5}\frac{c^5}{G}\left(\frac{r_s(M_h)}{a}\right)^5 f(e) \\ &= \frac{1}{5}\frac{c^5}{G}\left(\frac{r_s(M_r)}{a}\right)^2 \left(\frac{r_s(M)}{a}\right)^3 f(e), \end{aligned} \tag{8}$$

where

$$r_s(M_h) \equiv \frac{2GM_h}{c^2} \tag{9}$$

is the Schwarsczchild (horizon) radius of a non-spinning BH with mass M_h (hence the name "horizon" mass). The advantage of Equation (8) is that one can more straightforwardly track the dependence of $L_{\rm GW}$ on the separation between the two masses, a, i.e., $L_{\rm GW} \propto a^{-5}$. For $m_1 = m_2$, one has $M_h = 2^{1/5}m$, so that Equation (8) can be simplified to $L_{\rm GW} = (2/5)(c^5/G)(r_s(m)/a)^5$.

2.2. Electric Dipole Radiation Luminosity

When estimating the luminosity of the cCBC signal, Zhang (2016) considered the magnetic dipole radiation power by analogy with pulsars. Deng et al. (2018) noticed that for a cCBC system, electric dipole radiation is actually more significant. We treat the electric dipole radiation following Deng et al. (2018) in this subsection.

First, consider that only one member (m_2) is charged (with Q_2). The electric dipole radiation luminosity is

$$\begin{aligned} L_{\rm e,dip,2} &= \frac{2Q_2^2|\ddot{\boldsymbol{r}}_2|^2}{3c^3} \\ &= \frac{8}{3}\hat{q}_2^2\frac{G^3 m_1^2 m_2^2}{c^3 a^4} \\ &= \frac{1}{6}\frac{c^5}{G}\hat{q}_2^2\left(\frac{r_s(m_1)}{a}\right)^2\left(\frac{r_s(m_2)}{a}\right)^2, \end{aligned} \tag{10}$$

where $r_s(m_1)$ and $r_s(m_2)$ are the Schwarzschild radii for masses m_1 and m_2, respectively, and $|\ddot{\boldsymbol{r}}_2| = Gm_1/a^2$ is the amplitude of the acceleration of Mass 2. Noticing the symmetric format with respect to the two masses in Equation (10), it is straightforward to write down the general formula that both masses are charged:

$$L_{\rm e,dip} = \frac{1}{6}\frac{c^5}{G}(\hat{q}_1^2 + \hat{q}_2^2)\left(\frac{r_s(m_1)}{a}\right)^2\left(\frac{r_s(m_2)}{a}\right)^2. \tag{11}$$

There are two notes here. First, as the luminosity depends on $\hat{q}_1^2 + \hat{q}_2^2$, the power enhances no matter whether the two masses have the same or opposite charges. This is generally consistent with the numerical results of Liebling & Palenzuela (2016). Second, one may write down the ratio

$$\begin{aligned} \frac{L_{\rm e,dip}}{L_{\rm GW}} &= \frac{5}{6}(\hat{q}_1^2 + \hat{q}_2^2)\left(\frac{[r_s(m_1)]^2[r_s(m_2)]^2}{[r_s(M_h)]^4}\right)\left(\frac{a}{r_s(M_h)}\right) \\ &= \frac{5}{6}(\hat{q}_1^2 + \hat{q}_2^2)\left(\frac{M_r}{M}\right)^{2/5}\left(\frac{a}{r_s(M_h)}\right), \end{aligned} \tag{12}$$

which suggests that the electric dipole luminosity rises more slowly than the GW chirp signal.

One can also calculate the total dipole radiation energy through integrating luminosity over time. From an initial separation a to the final separation $a_{\min}$, one gets

$$\begin{aligned} E_{\rm e,dip} &= \int_a^{a_{\min}} da \frac{L_{\rm e,dip}}{\dot{a}} \\ &= \frac{5}{24} \ln\left(\frac{a}{a_{\min}}\right) (\hat{q}_1^2 + \hat{q}_2^2) M_r c^2, \end{aligned} \tag{13}$$

where we have used the orbital decay due to the GW loss (Maggiore 2008)

$$\dot{a} = -\frac{64}{5}\frac{G^3 M_r M^2}{c^5 a^3} f(e), \tag{14}$$

with $e = 0$ (and hence, $f(e) = 1$) adopted. As e also decreases with time for any elliptical orbits, it is reasonable to assume that the orbits are circular when a CBC occurs. Notice that Equation (13) does not converge as $a \to \infty$. However, because the power increases rapidly at the coalescence, one may mostly care about the energy release during the final orbits. The results are not sensitive to the actual values of a and $a_{\min}$ due to the logarithmic factor $\ln(a/a_{\min})$.

2.3. Magnetic Dipole Radiation Luminosity

For magnetic dipole radiation, we follow the same procedure of Zhang (2016). The magnetic dipole for the most general case is

$$\begin{aligned} \mu &= (\pi/c) I (a/2)^2, \\ &= \sqrt{GMa}\,(Q_1 + Q_2)/8c, \end{aligned} \tag{15}$$

where $I = (Q_1 + Q_2)/P$ is the current, and $P = 2\pi (GM)^{-1/2} a^{3/2}$ is the orbital period. The second derivative of μ reads

$$\ddot{\mu} = \frac{\sqrt{GM}}{16c}(Q_1 + Q_2)\left(-\frac{1}{2}a^{-3/2}\dot{a}^2 + a^{-1/2}\ddot{a}\right). \tag{16}$$

Again using Equation (14) with $f(e) = 1$, one gets $\ddot{a} = -(12288/25)(G^6 M_r^2 M^4 / a^7 c^{10})$. The magnetic dipole radiation power is then

$$\begin{aligned} L_{\rm B,dip} &= \frac{2\ddot{\mu}^2}{3c^3} \\ &= \frac{2^{17}\cdot 7^2}{3\cdot 5^4}\frac{c^5}{G}\left(\frac{\hat{q}_1 m_1 + \hat{q}_2 m_2}{M}\right)^2 \frac{G^{12} M_r^4 M^{11}}{c^{30} a^{15}} \\ &= \frac{196}{1875}\frac{c^5}{G}\left(\frac{\hat{q}_1 m_1 + \hat{q}_2 m_2}{M}\right)^2 \\ &\times \left(\frac{r_s(M_r)}{a}\right)^4 \left(\frac{r_s(M)}{a}\right)^{11}. \end{aligned} \tag{17}$$

When $m_1 = m_2$, this equation is reduced to Equation (7) of Zhang (2016).

One can also write down the ratios

$$\begin{aligned} \frac{L_{\rm B,dip}}{L_{\rm GW}} &= \frac{196}{375}\left(\frac{\hat{q}_1 m_1 + \hat{q}_2 m_2}{M}\right)^2 \\ &\times \left(\frac{r_s(M_r)}{a}\right)^2 \left(\frac{r_s(M)}{a}\right)^8; \end{aligned} \tag{18}$$

$$\begin{aligned} \frac{L_{\rm B,dip}}{L_{\rm e,dip}} &= \frac{392}{625}\frac{(\hat{q}_1 m_1 + \hat{q}_2 m_2)^2}{(\hat{q}_1^2 + \hat{q}_2^2) M^2} \\ &\times \left(\frac{r_s(M_r)}{a}\right)^2 \left(\frac{r_s(M)}{a}\right)^9, \end{aligned} \tag{19}$$

which suggest that the magnetic dipole radiation power rises much faster than both the GW power and the electric dipole radiation power. At large separations ($a \gg r_s M_h$), this term is negligibly small. However, at the merger time, $a \sim r_s(m_1) + r_s(m_2) = r_s(M)$ for BH–BH mergers, $L_{\rm B,dip}$ becomes comparable to $L_{\rm e,dip}$.

The total magnetic dipole radiation energy can be obtained as

$$\begin{aligned} E_{\rm B,dip} &= \int_a^{a_{\min}} da \frac{L_{\rm B,dip}}{\dot{a}} \\ &= \frac{49}{4125}(M_r c^2)\left(\frac{\hat{q}_1 m_1 + \hat{q}_2 m_2}{M}\right)^2 \\ &\times \left(\frac{r_s(M_r)}{a_{\min}}\right)^2 \left(\frac{r_s(M)}{a_{\min}}\right)^9, \end{aligned} \tag{20}$$

which depends on $a_{\min}$. For BH–BH mergers, one has $a_{\min} \sim r_s(m_1) + r_s(m_2) = r_s(M)$, so that

$$E_{\rm B,dip} = \frac{49}{4125}(M_r c^2)\left(\frac{\hat{q}_1 m_1 + \hat{q}_2 m_2}{M}\right)^2 \left(\frac{M_r}{M}\right)^2. \tag{21}$$

The ratio between the two total energy components is

$$\frac{E_{\rm B,dip}}{E_{\rm e,dip}} = \frac{392}{6875}\frac{(\hat{q}_1^2 m_1 + \hat{q}_2 m_2)^2}{(\hat{q}_1^2 + \hat{q}_2^2) M^2 \ln(a/a_{\min})}\left(\frac{M_r}{M}\right)^2. \tag{22}$$

One can see that usually $E_{\rm B,dip} \ll E_{\rm e,dip}$. For $m_1 = m_2$, $\hat{q}_1 = \hat{q}_2$, this ratio is $\sim 1.8 \times 10^{-3} \ln^{-1}(a/a_{\min})$.

2.4. Radiation Signature

The discussion so far does not specify the form of EM radiation that these systems emit. Both electric and magnetic dipole radiations have the frequency of the orbital frequency of the system, which falls around the kHz range for CBCs. This frequency is below the typical plasma frequency $\omega_p = (4\pi n e^2/m)^{1/2} = (5.64 \times 10^4\ {\rm Hz})\, n^{1/2}$ for a typical interstellar medium with $n \sim 1\ {\rm cm}^{-3}$, so the dipole radiations themselves cannot propagate. In reality, the radiation energy is advected in the form of an outgoing Poynting flux dominated outflow. Particle acceleration and subsequent radiation would occur within the outflow with the expense of the Poynting flux energy, so that broadband radiation (from radio to γ-rays) is possible. This is certainly the case for spindown-powered pulsars, whose magnetic dipole radiation is released in the form of a pulsar wind and various forms of radiations (e.g., for the Crab pulsar, coherent radio emission, nonthermal γ-ray, and X-ray emission due to photon-pair cascade from the magnetosphere or current sheet outside the light cylinder, and broadband pulsar wind nebula radiation in a much larger scale). Similar processes may happen for cCBCs, at least for magnetic dipole radiation, but possibly for electric dipole radiation as well. Zhang (2016) discussed the possible mechanisms of converting the magnetic dipole radiation to the magnetospheric coherent radio emission (to power an FRB) and the internal Poynting-flux-dissipation-powered γ-ray emission (to power a weak GRB).

3. EM Counterpart of Plunging BH–NS Mergers

For plunging BH–NS mergers, i.e., the mass ratio is required to be $q < 0.2$ (Shibata et al. 2009). The NS is swallowed by the BH as a whole, so that no matter-related EM counterparts (short GRB and kilonova) are expected.

As has been well known in pulsar theories (e.g., Michel 1982), rotating, magnetized NSs are globally charged. This is because the requirement that the plasma co-rotates with the NS (both interior and exterior to the NS surface) is that charges are spatially separated to maintain a certain charge density distribution (Goldreich & Julian 1969; Ruderman & Sutherland 1975). When integrating the Goldreich–Julian (GJ) spatial charge density distribution ($\rho_{\rm GJ} \sim -(\boldsymbol{\Omega} \cdot \boldsymbol{B})/2\pi c$) over a volume, one obtains a net charge. Such a net charge has been observed from particle-in-cell numerical simulations of force-free pulsars (Pétri 2012). In general, the net charge decreases with the inclination angle of the NS (approaching zero for perpendicular rotators) and the volume integrated (in the scale comparable to light cylinder and beyond); see Figures 12 and 13 of Pétri (2012). For simplicity, we consider a rotator with the spin axis and magnetic axis anti-parallel ($\boldsymbol{\Omega} \cdot \boldsymbol{B} < 0$). Because the coalescence is relevant only the near-zone magnetosphere, we do not consider the depletion of the net charge at large distances. The charge contained in the magnetosphere can be obtained by integrating the GJ density from the NS surface R to an arbitrarily large radius r (e.g., the light cylinder). Assuming a magnetic dipole field, i.e., $B_r(r, \theta) = (\mu_*/r^3) \cdot (2\cos\theta)$ and $B_\theta(r, \theta) = (\mu_*/r^3) \cdot \sin\theta$ (where μ_* is the dipole moment of the NS, in contrast with μ, which is the dipole moment of the binary system), the integrated charge between radii R and r is

$$
\begin{aligned}
Q_{\rm mag} &= 4\pi \int_R^r \int_0^{\pi/2} -\frac{\boldsymbol{\Omega} \cdot \mathbf{B}}{2\pi c} r^2 d\theta dr \\
&= \int_R^r \frac{1}{r} dr \int_0^{\pi/2} \frac{2\Omega\mu_*}{c} \cos\theta \sqrt{1 + 3\cos^2\theta}\, d\theta \\
&= \ln\left(\frac{r}{R}\right) \frac{\Omega\mu_*}{c} \left(1 + \frac{4\sqrt{3}}{9}\pi\right), \\
&= \ln\left(\frac{r}{R}\right) \frac{\Omega B_p R^3}{c} \left(\frac{1}{2} + \frac{2\sqrt{3}}{9}\pi\right),
\end{aligned} \tag{23}
$$

which is insensitive to R and r. The charge contained within the NS can be calculated similarly. For a dipole configuration inside the NS, one can adopt the same formula (Equation (23)) with the inner radius set at where the dipole approximation is broken. For a uniformly magnetized NS, the NS charge is $Q_{\rm NS} = (\Omega B)/(2\pi c) \cdot (4\pi/3) R^3 = (2/3)\Omega B_p R^3$. In the following, we estimate the total charge $Q_{\rm tot} = Q_{\rm mag} + Q_{\rm NS} \sim (3\Omega B_p R^3/c)(\cos\alpha)$, where the dependence on the inclination angle α (Pétri 2012) has been included. For the typical NS mass $M = (1.4\, M_\odot) M_{1.4}$, this gives a dimensionless NS charge

$$
\hat{q}_{\rm NS} \simeq \frac{3\Omega B_p R^3}{2c\sqrt{G}M} \cos\alpha = (4.4 \times 10^{-7}) B_{13} P_{-2}^{-1} R_6^3 M_{1.4}^{-1} \cos\alpha. \tag{24}
$$

For the Crab pulsar, one has $B_{13} = 0.8$ and $P_{-2} = 3.3$, which gives $\hat{q}_{\rm Crab} = 1.1 \times 10^{-7} \cos\alpha$.

For a plunging BH–NS merger system with $q = 0.2$, assuming that the BH charge is much smaller than the NS charge, one can calculate maximum luminosities and total energies of both electric and magnetic dipole radiations using Equations (11), (13), (17), and (21). Notice that at the merger (when the NS touches the BH), one has $a_{\rm min} = r_s(m_1) + 2.4 r_s(m_2)$ (the radius of an NS is about 2.4 times larger than its Schwarzschild radius), we finally get ($a/a_{\rm min} = 10$ is adopted)

$$
\begin{aligned}
L_{\rm e,dip,max} &= (5.0 \times 10^{42}\ {\rm erg\ s^{-1}})\ \hat{q}_{-7}^2, \\
E_{\rm e,dip} &= (1.0 \times 10^{40}\ {\rm erg})\ \hat{q}_{-7}^2, \\
L_{\rm B,dip,max} &= (1.7 \times 10^{38}\ {\rm erg\ s^{-1}})\ \hat{q}_{-7}^2, \\
E_{\rm B,dip} &= (1.2 \times 10^{34}\ {\rm erg})\ \hat{q}_{-7}^2.
\end{aligned} \tag{25}
$$

One can see that for $\hat{q} \sim 10^{-7}$, the electric dipole luminosity/energy reaches that of an FRB assuming isotropic radiation (e.g., Zhang 2018 for a detailed calculation of FRB energetics). Whether such an EM pulse can be emitted in the GHz frequency range to power an FRB is subject to further studies. In any case, the EM signal is brief and has the right luminosity and energy to potentially power a non-repeating FRB.

4. EMRI Systems

Besides plunging BH–NS events, another type of system to observe cCBC signals is EMRI systems (e.g., Amaro-Seoane et al. 2007), which are possible targets for space GW detectors such as Laser Interferometer Space Antenna (LISA). For such systems, as $q = m_2/m_1 \ll 1$, one has $M_r \sim m_2$, $M \sim m_1$, and $a_{\rm min} \sim r_s(m_1)$. Equations (11), (13), (17), and (21) can be written as

$$
L_{\rm e,dip,EMRI} = (6.1 \times 10^{36}\ {\rm erg\ s^{-1}}) \hat{q}_{-7}^2 q_{-4}^2 \left(\frac{a_{\rm min}}{a}\right)^4; \tag{26}
$$

$$
E_{\rm e,dip,EMRI} = (5.2 \times 10^{39}\ {\rm erg}) \hat{q}_{-7}^2 \ln\left(\frac{a}{a_{\rm min}}\right); \tag{27}
$$

$$
L_{\rm B,dip,EMRI} = (3.8 \times 10^{20}\ {\rm erg\ s^{-1}}) \hat{q}_{-7}^2 q_{-4}^6 \left(\frac{a_{\rm min}}{a}\right)^{15}; \tag{28}
$$

$$
E_{\rm B,dip,EMRI} = (3.0 \times 10^{22}\ {\rm erg}) \hat{q}_{-7}^2 q_{-4}^4 \left(\frac{a_{\rm min}}{a}\right)^{11}. \tag{29}
$$

One can see that magnetic dipole radiation is too weak for any observational interest. If electric dipole radiation can be converted to detectable signals (e.g., radio emission), there might be a possibility of detection by continuing to observe targeted sources over a long period of time. For example, for $q = 10^{-4}$ and $a = 2a_{\rm min}$, the electric dipole radiation luminosity is $\sim 3.8 \times 10^{35}\ {\rm erg\ s^{-1}}$. Even if the luminosity is much lower than that of stellar-mass BH–NS systems, these systems are long lived, and the signal becomes stronger after a long-time integration. If LISA identifies an EMRI source in the sky with the NS very close to the event horizon of the massive BH, long-term radio monitoring is recommended to detect the possible cCBC signal associated with the system.

5. Summary and Discussion

Continuing from a previous investigation (Zhang 2016), in this Letter we develop a more general theory for CBC systems with at least one member charged and study the cCBC EM signal in great detail. The current treatment allows different masses and amounts of charge for the two merging members, and the luminosities and energies due to both electric and magnetic dipole radiations are calculated. The most useful expressions are Equations (11), (13), (17), and (21). In general, because $\hat{q}$ of astrophysical objects is typically $\ll 1$, the cCBC signal is not expected to be very bright.

Even though it is still an open question whether significantly charged Kerr–Newmann BHs can survive for a long enough time, it is well known that NSs are globally charged. Therefore, the cCBC signal should present in NS mergers. This signal is likely non-detectable in NS–NS mergers or BH–NS mergers with $q > 0.2$, when plenty of matter is present outside of the horizon of the post-merger BH. The plunging BH–NS mergers with $q < 0.2$ are ideal systems to cleanly observe the cCBC signal. Our estimate suggests that in order to have the cCBC signal reaching a detectable level (e.g., that of an FRB), the NS needs to be young with a relatively strong magnetic field, e.g., Crab-like. This may be possible in young star clusters where BH–NS binaries may form in tight orbits so that the merger occurs within a short period of time when the pulsar is still young. As the charge of the NS $Q \propto \Omega B_p$, the cCBC signal, if detected, can be used to infer the NS parameter (ΩB_p) before falling into the BH.

In addition to the cCBC signal, the distortion of the NS magnetosphere itself near the end of CBC may trigger more magnetic dissipation in the NS magnetosphere, the strength of which depends on the available magnetic energy. For an order of magnitude estimate, the total dissipated energy would be $\sim(B^2/8\pi)R^3 \sim (4 \times 10^{40}\ \mathrm{erg})\ B_{12}^2 R_6^3$. One can see that in order to produce an FRB-like event, one needs to dissipate a strong enough magnetic field in a large enough volume. For a favorable magnetic configuration, the requirement for pulsar parameters may be less demanding.

Dai (2019) recently proposed another interesting cCBC signal related to BH-NS mergers. In his scenario, if the BH is rapidly spinning, it will progressively gain charge from the NS magnetosphere through the "Wald mechanism" (Wald 1974) as the two merging objects approach each other. His signal may become stronger than the signal discussed here if the BH spins fast enough.

Finally, EMRIs are another type of target to search for the cCBC signal. The systems with a relatively small mass ratio q (within the EMRI category) and the ratio $a/a_{\min}$ close to unity are favorable targets to search for such a signal in the era of space GW astronomy.

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References

Abbott, B. P., Abbott, R., Abbott, T. D., et al. 2017a, ApJL, 848, L13
Abbott, B. P., Abbott, R., Abbott, T. D., et al. 2017b, PhRvL, 119, 161101
Amaro-Seoane, P., Gair, J. R., Freitag, M., et al. 2007, CQGra, 24, R113
Barnes, J., & Kasen, D. 2013, ApJ, 775, 18
Bartos, I., Brady, P., & Márka, S. 2013, CQGra, 30, 123001
Connaughton, V., Burns, E., Goldstein, A., et al. 2016, ApJL, 826, L6
Connaughton, V., Burns, E., Goldstein, A., et al. 2018, ApJL, 853, L9
Coulter, D. A., Foley, R. J., Kilpatrick, C. D., et al. 2017, Sci, 358, 1556
Dai, L., McKinney, J. C., & Miller, M. C. 2017, MNRAS, 470, L92
Dai, Z. G. 2019, ApJL, in press (arXiv:1902.07939)
Deng, C.-M., Cai, Y., Wu, X.-F., & Liang, E.-W. 2018, PhRvD, 98, 123016
Eichler, D., Livio, M., Piran, T., & Schramm, D. N. 1989, Natur, 340, 126
Fraschetti, F. 2018, JCAP, 4, 054
Gao, H., Ding, X., Wu, X.-F., Zhang, B., & Dai, Z.-G. 2013, ApJ, 771, 86
Goldreich, P., & Julian, W. H. 1969, ApJ, 157, 869
Goldstein, A., Veres, P., Burns, E., et al. 2017, ApJL, 848, L14
Greiner, J., Burgess, J. M., Savchenko, V., & Yu, H.-F. 2016, ApJL, 827, L38
Hallinan, G., Corsi, A., Mooley, K. P., et al. 2017, Sci, 358, 1579
Kasen, D., Metzger, B., Barnes, J., Quataert, E., & Ramirez-Ruiz, E. 2017, Natur, 551, 80
Kimura, S. S., Takahashi, S. Z., & Toma, K. 2017, MNRAS, 465, 4406
Levin, J., D'Orazio, D. J., & Garcia-Saenz, S. 2018, PhRvD, 98, 123002
Li, L.-X., & Paczyński, B. 1998, ApJL, 507, L59
Liebling, S. L., & Palenzuela, C. 2016, PhRvD, 94, 064046
Liu, T., Romero, G. E., Liu, M.-L., & Li, A. 2016, ApJ, 826, 82
Loeb, A. 2016, ApJL, 819, L21
Lorimer, D. R., Bailes, M., McLaughlin, M. A., Narkevic, D. J., & Crawford, F. 2007, Sci, 318, 777
Lyman, J. D., Lamb, G. P., Levan, A. J., et al. 2018, NatAs, 2, 751
Maggiore, M. 2008, Gravitational Waves. Volume 1: Theory and Experiments (Oxford: Oxford Univ. Press)
Margutti, R., Alexander, K. D., Xie, X., et al. 2018, ApJL, 856, L18
Mészáros, P., & Rees, M. J. 1992, MNRAS, 257, 29P
Metzger, B. D., Martínez-Pinedo, G., Darbha, S., et al. 2010, MNRAS, 406, 2650
Metzger, B. D., & Piro, A. L. 2014, MNRAS, 439, 3916
Michel, F. C. 1982, RvMP, 54, 1
Narayan, R., Paczynski, B., & Piran, T. 1992, ApJL, 395, L83
Perna, R., Lazzati, D., & Giacomazzo, B. 2016, ApJL, 821, L18
Pétri, J. 2012, MNRAS, 424, 605
Piro, L., Troja, E., Zhang, B., et al. 2019, MNRAS, 483, 1912
Ruderman, M. A., & Sutherland, P. G. 1975, ApJ, 196, 51
Shibata, M., Kyutoku, K., Yamamoto, T., & Taniguchi, K. 2009, PhRvD, 79, 044030
Thornton, D., Stappers, B., Bailes, M., et al. 2013, Sci, 341, 53
Troja, E., Piro, L., van Eerten, H., et al. 2017, Natur, 551, 71
Villar, V. A., Guillochon, J., Berger, E., et al. 2017, ApJL, 851, L21
Wald, R. M. 1974, PhRvD, 10, 1680
Yu, Y.-W., Zhang, B., & Gao, H. 2013, ApJL, 776, L40
Zhang, B. 2013, ApJL, 763, L22
Zhang, B. 2016, ApJL, 827, L31
Zhang, B. 2018, ApJL, 867, L21
Zhang, B.-B., Zhang, B., Sun, H., et al. 2018, NatCo, 9, 447