MERGERS OF CHARGED BLACK HOLES: GRAVITATIONAL-WAVE EVENTS, SHORT GAMMA-RAY BURSTS, AND FAST RADIO BURSTS

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Received 2016 March 22; revised 2016 July 12; accepted 2016 July 12; published 2016 August 16

ABSTRACT
The discoveries of GW150914, GW151226, and LVT151012 suggest that double black hole (BH–BH) mergers are common in the universe. If at least one of the two merging black holes (BHs) carries a certain amount of charge, possibly retained by a rotating magnetosphere, the inspiral of a BH–BH system would drive a global magnetic dipole normal to the orbital plane. The rapidly evolving magnetic moment during the merging process would drive a Poynting flux with an increasing wind power. The magnetospheric activities during the final phase of the merger would make a fast radio burst (FRB) if the BH charge can be as large as a factor of \( \hat{q} \sim (10^{-5} – 10^{-8}) \) of the critical charge \( q_c \) of the BH. At large radii, dissipation of the Poynting flux energy in the outflow would power a short-duration high-energy transient, which would appear as a detectable short-duration gamma-ray burst (GRB) if the charge can be as large as \( \hat{q} \sim (10^{-5} – 10^{-9}) \). The putative short GRB coincident with GW150914 recorded by Fermi GBM may be interpreted with this model. Future joint GW/GRB/FRB searches would lead to a measurement or place a constraint on the charges carried by isolate BHs.

Key words: radiation mechanisms: non-thermal – stars: black holes

1. INTRODUCTION
Black holes (BHs) are uniquely described with three parameters, mass \( M \), angular momentum \( J \), and charge \( Q \). Whereas the first two parameters have been measured with various observations for both stellar-mass and super-massive BHs, it has been widely believed that the \( Q \) parameter must be very small. However, no measured value or upper limit of \( Q \) have been reported for any BH.

Recently, the Laser Interferometer Gravitational-wave Observatory team announced the ground-breaking discovery of the first gravitational-wave (GW) source, GW150914, which is a double black hole (BH–BH) merger with two BH masses \( 36.4_{-2}^{+5} M_\odot \) and \( 29.4_{-4}^{+2} M_\odot \), respectively (Abbott et al. 2016a). Two other BH–BH merger events (GW151226 and LVT151012) were later announced (Abbott et al. 2016b). The inferred event rate density of BH–BH mergers is \( \sim (9–240) \text{ Gpc}^{-3} \text{yr}^{-1} \) (Abbott et al. 2016c). Intriguingly, the Fermi GBM team reported a 1 s long, putative, weak gamma-ray burst (GRB) 0.4 s after the GW event was detected (Connaughton et al. 2016; but see Greiner et al. 2016; Xiong 2016). This is surprising, since unlike NS–NS and NS–BH mergers that can form BH–torus systems and produce short GRBs through accretion (Paczynski 1986; Eichler et al. 1989; Paczynski 1991; Mészáros & Rees 1992; Narayan et al. 1992; Rezzolla et al. 2011), BH–BH mergers are not expected to have surrounding materials with a high enough density to power a short-duration GRB via accretion.

On the other hand, fast radio bursts (FRBs) are mysterious millisecond-duration radio transients (Lorimer et al. 2007; Thornton et al. 2013). Recent observations suggest that at least some FRBs are likely at cosmological distances (e.g., Keane et al. 2016). Their physical origins, however, remain unknown.

Here, we show that if at least one BH in the two merging BHs carries a certain amount of charge, the inspiral of the BH–BH system would induce a global magnetic dipole normal to the orbital plane. The rapid evolution of the magnetic moment would drive a Poynting flux with an increasing wind power, which may give rise to an FRB and even a short-duration GRB depending on the value of the charge.

2. ELECTRODYNAMICS OF CHARGED BH MERGER SYSTEM
For a charged BH, one can define the Schwarzschild radius and the Reissner–Nordström radius
\[
r_s = \frac{2GM}{c^2}, \quad r_0 = \frac{\sqrt{GQ}}{c^2},
\]
where \( M \) and \( Q \) are the mass and charge of the BH, respectively; \( G \) and \( c \) are the gravitational constant and speed of light, respectively; and the electrostatic cgs units have been used. By equating \( r_s \) and \( r_0 \), one may define a characteristic charge
\[
Q_c \equiv 2\sqrt{G}M = (1.0 \times 10^{31} \text{e.s.u.}) \left( \frac{M}{10 M_\odot} \right),
\]
which is \( (3.3 \times 10^{21} \text{C}) (M/10 M_\odot) \) in S.I. units. The charge of this magnitude would significantly modify the spacetime geometry with a magnitude similar to \( M \). We consider a BH with charge
\[
Q = \hat{q} Q_c,
\]
with the dimensionless parameter \( \hat{q} \ll 1 \). For simplicity, in the following, we consider two identical BHs with the same \( M \) and \( Q \).

As the two BHs spiral in, a circular current loop forms, this magnitude would significantly modify the spacetime geometry with a magnitude similar to \( M \). We consider a BH with charge
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\[
Q = \hat{q} Q_c,
\]
where $I = 2Q/P$ is the current, and

$$P = \frac{2\pi}{\sqrt{2GM}} a^{3/2} = \frac{8\sqrt{2} \pi GM}{c^3} \dot{a}^{3/2}$$

$$= (1.7 \text{ ms}) \left( \frac{M}{10 M_\odot} \right) \dot{a}^{3/2} \tag{5}$$

is the Keplerian orbital period, $a = \dot{a}(2r)$ is the separation between the two BHs, and $\dot{a}$ is the distance normalized to $2r$. Notice that at the coalescence of the two BHs, $\dot{a} = 1$ for two Schwarzschild BHs, but $\dot{a}$ can be as small as 0.5 for extreme Kerr BHs. For comparison, a magnetar with a surface magnetic field $B_p \sim 10^{15} G$ and radius $R_{\text{NS}} \sim 10^6$ cm has a magnetic dipole $\mu_{\text{mag}} \sim B_p R_{\text{NS}}^3 = (10^{33} \text{ G cm}^3) B_{p,15} R_{\text{NS,6}}^3$.

The orbital decay rate due to GW radiation can be generally written as $da/dt = -(64/5)GM a_m^2 /c^2 a^2 (1 - e^2/2)(1 + (73/24)e^2 + (37/96)e^4)$, where $a = M_1 M_2 /M_{\text{tot}}$ is the chirp mass and $M_{\text{tot}} = M_1 + M_2$ is the total mass of the system. Assuming $M_1 = M_2$ for simplicity and adopting $e = 0$, which is valid before the coalescence, one gets

$$\frac{da}{dt} = -\frac{2}{5} \frac{c}{\dot{a}}. \tag{6}$$

The rapid evolution of the orbital separation before the coalescence leads to a rapid change of the magnetic flux, and hence a Poynting flux with increasing power. A full description of the electrodynamics of the system requires numerically solving Einstein equations with electrodynamics. To an order of magnitude analysis, one may estimate the Poynting-flux wind power

$$L_w \sim \dot{E}^2 /R, \tag{10}$$

where $R$ is the resistance of the magnetosphere, which may be taken as $c^{-1}$ for a conductive magnetosphere. This gives

$$L_w \sim \dot{E}^2 = \frac{\pi^2 GM^2}{50} \epsilon^2 \gamma^2 \dot{a}^{-7} \simeq \frac{\pi^2}{200} \epsilon^2 \gamma^2 \dot{a}^{-7}, \tag{11}$$

where $\dot{E} = r/2r_t$ is the normalized wind-launching radius. Notice that Equation (11) has the same scaling $\sim (c^5/G) \dot{q}^2$ as Equation (7), even though the dependence on $\dot{a}$ may be different (pending on how $\dot{r}$ depends on $\dot{a}$). In the following, for simplicity, we apply the vacuum formula Equation (7) to perform related estimates.

The wind power is very sensitive to $\dot{a}$ and increases rapidly as the orbital separation shrinks. The highest power happens right before the final merger so that such a merger system is a plausible engine for an FRB and possibly a short-duration $\gamma$-ray burst.\(^2\)

One may estimate the timescale for the orbital separation to shrink from $\dot{a} = 1.5$ to $\dot{a} = 1$, during which $L_w$ increases by a factor of $\sim 440$. This is

$$\tau_{1.5} \lesssim \frac{P}{|\dot{P}|} = \frac{20GM}{3} \dot{a}^{4} \dot{a} \simeq (1.7 \text{ ms}) \left( \frac{M}{10 M_\odot} \right) \dot{a}^{4}, \tag{12}$$

where $\dot{P} \simeq -(192\pi/5c^5)(2\pi G/P)^{5/3} M^2 M_{\text{tot}}^{-1/3} = (6\sqrt{2} \pi /5) \dot{a}^{-5/2}$ is the orbital decay rate for GW radiation (Taylor & Weissberg 1989).

It would be informative to compare the Poynting flux power proposed in this Letter (Equation (7)) with some other Poynting flux powers proposed in the literature. Two relevant ones are the general-relativity-induced Poynting flux power when a BH moves in a constant magnetic field $B_0$ (Lyutikov 2011a)\(^3\) and a Poynting flux power due to the interaction between the magnetospheres of two BHs (Lyutikov 2011b).\(^4\) Expressing Equations (1) and (4) in Lyutikov (2011b) in terms of $\dot{q}$ using Equation (13) below, we find that these two powers are both of the order of $-\left(\dot{R}_{\text{NS}}/a\right)^2 \dot{a}^2 L_w$, where $R_{\text{NS}} = c/\Omega_{\text{NS}}$ is the light cylinder radius of the BHs. Notice the strong dependence on $\dot{a}$. These powers are negligibly small compared with $L_w$ when $\dot{a}$ becomes smaller than unity.

\(^2\) After the submission of this Letter, Liu et al. (2016) proposed an alternative mechanism to produce FRBs from BH–BH merger systems through triggering an instability in the Kerr–Newman BH magnetospheres.

\(^3\) In a dynamically evolving system, the assumption of constant $B_0$ is no longer valid so that more detailed modeling is needed to perform a more accurate comparison between this power and $L_w$.

\(^4\) This power does not exist if only one BH carries a magnetosphere.
3. ON THE CHARGE OF BHs

It is well known that a rotating point magnetic dipole carries a net charge (Cohen et al. 1975; Michel 1982). In the physical model of pulsars, the difficulty was not how to make a charged neutron star, but rather how to designate a return current to make a neutron star neutral (which is not necessary in pulsar emission models; Michel 1982). We assume that the charged BHs in our model each possess a magnetosphere with a dipole configuration. The magnetosphere may be attained in the not-too-distant past when the BH went through a magnetically arrested accretion phase (e.g., Tchekhovskoy et al. 2011) and the BH is still undergoing slow “balding” (Lyutikov & McKinney 2011). Alternatively, the magnetosphere may be maintained by a debris disk that is circulating the BH at the time of coalescence (e.g., Li et al. 2016; Perna et al. 2016). The charge maintained by an astrophysical rotating dipolar magnetosphere is approximately

\[ Q \sim \frac{\Omega_0 \mu_0}{3c}, \]

(13)

where \( \mu_0 \) (to be differentiated from \( \mu \) in Equation (4)) is the magnetic moment of the BH dipole, and \( \Omega_0 \) is the angular velocity of the BH magnetosphere. This may be derived according to the Gauss’s law for a point dipole (Michel 1982) or through a volume integration of a Goldreich–Julian magnetosphere.

According to Equations (2) and (13), the rotating magnetic point dipole of individual BHs with dimensionless charge \( \hat{q} \) should satisfy

\[ \mu_0 \Omega_0 \sim (9 \times 10^{36} \text{ G cm}^3 \text{s}^{-1}) \left( \frac{M}{10 M_\odot} \right) \hat{q}^{1.5}. \]

(14)

For comparison, a millisecond pulsar has \( \mu_0 \Omega_0 \sim 10^{37} \text{ G cm}^3 \text{s}^{-1} \).

The spin-down luminosity of individual BHs with magnetic dipoles may be estimated as \( L_\nu \sim (2 \mu_0^2 \Omega_0^2)/(3c^3) \). This gives

\[ \frac{L_\nu}{L_w} \sim \left( \frac{120 \tau_{\text{sc}}}{R_{bcu}} \right)^2 \hat{a}^{15} \sim 0.4 \left( \frac{r_c}{R_{bcu}} \right)^2 \left( \frac{\hat{a}}{0.5} \right)^{15}. \]

(15)

One can see that even though \( L_\nu \gg L_w \) when \( \hat{a} \gg 1 \), at coalescence (\( \hat{a} \ll 1 \)) \( L_\nu \) becomes smaller than \( L_w \). In the slow-balding scenario of Lyutikov & McKinney (2011), the field would evolve into a monopole configuration. In this case, one may estimate \( L_\nu \sim (\Omega_0 \mu_0 / r_c)^2 / c \sim (c^3 / G) \hat{q}^2 \). This gives

\[ \frac{L_\nu}{L_w} \sim 2400 \hat{a}^{15} \sim 0.07 \left( \frac{\hat{a}}{0.5} \right)^{15}. \]

Again thanks to the strong dependence of \( L_w \) on \( \hat{a} \), \( L_\nu \) becomes negligibly small compared with \( L_w \) at \( \hat{a} < 1 \).

4. RADIO AND GAMMA-RAY EMISSION

In this model, radio emission may be produced in the inner magnetosphere through a coherent “bunching” curvature radiation mechanism by the pairs streaming out from the magnetosphere, similar to the case of radio pulsars. The timescale (Equation (12)) sets an upper limit on the duration of an FRB. To reproduce a typical FRB luminosity \( L_{\text{FRB}} \sim 10^{41} \text{ erg s}^{-1} \), the requirement of \( L_w > L_{\text{FRB}} \) (from Equation (7)) gives \( \hat{q} > 3 \times 10^{-3} \) for \( \hat{a} = 1 \) and \( \hat{q} > 2 \times 10^{-10} \) for \( \hat{a} = 0.5 \).

The magnetic field configuration of the dynamical magnetosphere is complicated. For simplicity, we adopt a dipole field as an order of magnitude estimate. Right before the coalescence, one has \( a = (4GM/c^2)\hat{a} = (1.8 \times 10^3 \text{ cm})(M/30 M_\odot)\hat{a} \) and \( \hat{a} \gg 1 \). For a dipole field line \( r = r_c \sin^2 \theta \), one may take \( r_c \sim a \) right before the coalescence (which implies a nearly isotropic emission beam). Noticing that the curvature radius \( \rho \sim (0.3-0.6)r_c \) in a wide range of \( r \), one may approximate \( \rho \sim 0.45r_c \sim (8 \times 10^6 \text{ cm})(M/30 M_\odot)\hat{a} \). The typical curvature radiation frequency of the pairs is

\[ \nu = \frac{3}{4\pi} \frac{c}{\rho} \gamma_\epsilon \gamma_e \sim (0.9 \times 10^9 \text{ Hz}) \hat{a}^{-1} \left( \frac{M}{10 M_\odot} \right)^{-1} \gamma_\epsilon \gamma_e e^{-}, \]

(17)

where the Lorentz factor of the pairs \( \gamma_\epsilon \) is normalized to 100, the nominal Lorentz factor value of pairs from a pulsar polar cap cascade (e.g., Zhang & Harding 2000). This frequency is the typical frequency of the observed FRBs. The curvature radiation emission power of an electron is \( P_e = \frac{2}{3} \epsilon \gamma_\epsilon \gamma_e \sim (7.2 \times 10^{-15} \text{ erg s}^{-1}) \hat{a}^{-2} (M/10 M_\odot)^{-2} \gamma_\epsilon \gamma_e e^{-} \). For the bunching coherent mechanism (Ruderman & Sutherland 1975), the total emission power is \( P = N_{\text{bunch}} N_e^2 P_e \), where \( N_e \) is the number of electrons in each bunch, \( N_{\text{bunch}} \) is the number of bunches, with the total number of electrons defined by \( N_{\text{tot}} = N_{\text{bunch}} N_e \). The minimum number of electrons that are needed to reproduce the typical luminosity of an FRB, \( L_{\text{FRB}} = 10^{41} \text{ erg s}^{-1} L_{\text{FRB,41}} \), can be derived by assuming that \( N_{\text{bunch}} = 1 \) and \( N_{\text{tot}} = N_e \) so that \( N_{\text{tot,min}} = (L_{\text{FRB}}/P_e)^{1/2} \sim 3.7 \times 10^{27} \hat{a}(M/10 M_\odot)\gamma_\epsilon \gamma_e L_{\text{FRB,41}} \). The total number of emitting electrons in the magnetosphere may be estimated as \( N_{\text{tot}} \sim \epsilon \gamma_e \sim (2.1 \times 10^{31} \text{ s}^{-1})q_{\gamma_e} \), which is \( \gg N_{\text{tot,min}} \) even if \( \hat{q} \) is normalized to \( 10^{-9} \). This suggests that energetically the bunching mechanism is able to power an FRB in such a transient magnetosphere.

The pair cascade process only converts a small fraction of the wind energy into radio emission. The dominant energy component in the outflow would be in the form of a Poynting flux. The EM energy is entrained in the outflow and would be dissipated at a large radius through magnetic reconnection triggered by internal collision (Zhang & Yan 2011) or current instabilities (Lyutikov & Blandford 2003). Assuming that GWs travel with the speed of light,\(^5\) the FRB is essentially simultaneous with the GW chirp signal, but the gamma-ray emission would be slightly delayed with respect to the GW chirp signal due to the slightly smaller speed of the Poynting flux with respect to the speed of light. Suppose that the GRB emission starts at radius \( R_1 \) with Lorentz factor \( \Gamma_1 \) and ends at radius \( R_2 \) with Lorentz factor \( \Gamma_2 \), one may define

\[ t_1 = \frac{R_1}{2 \Gamma_1 c}, \quad t_2 = \frac{R_2}{2 \Gamma_2 c}. \]

(18)

Several observational timescales can be estimated as follows:

1. The delay time between the onset of the GRB and the final GW chirp signal is

\[ \Delta t_{\text{GRB}} \sim (t_1 - t_2)(1+z). \]

(19)

\(^5\) The GW150914 indeed leads the putative associated GRB by 0.4 s (Connahhott et al. 2016). This would give the tightest constraint on Einstein’s Equivalent Principle to date (Wu et al. 2016).
2. The rising timescale of the GRB is defined by
\[ t_r \sim \max(\eta_s, t_2 - t_1)(1 + z). \]  
(20)

3. The decay timescale of the GRB is defined by
\[ t_d \sim t_2(1 + z). \]  
(21)

4. The total duration of the GRB is
\[ \tau \equiv t_r + t_d. \]  
(22)

5. GW150914 AND THE POSSIBLE ASSOCIATED GRB

Connaughton et al. (2016) reported a weak, hard X-ray transient that was potentially associated with GW150914. The false-alarm probability is 0.0022, and the poorly constrained localization is consistent with that of GW150914. The putative GRB has a duration \( \tau \sim 1 \) s and was delayed with respect to the GW signal by \( \Delta t_{\text{GRB}} \sim 0.4 \) s. Assuming the redshift of GW150914 (Abbott et al. 2016a), \( z = 0.09 \pm 0.03 \), the 1 keV–10 MeV luminosity of the putative GRB is \( 1.8 \times 10^{39} \) erg s\(^{-1}\).

The properties of this putative short GRB may be interpreted by our model. According to Equation (7), one can estimate the required charge of the BHs as
\[ \dot{q}_{-4} \approx 3.5\dot{\alpha}^{15/2}/\eta_\gamma^{-1/2} \approx 0.02\left(\frac{\dot{q}}{0.5}\right)^{15/2}/\eta_\gamma^{-1/2}, \]  
(23)

where \( \eta_\gamma \equiv L_\gamma/L_w \) is the radiative efficiency of the GRB, which ranges in (0.1–1) for known GRBs (Zhang et al. 2007). According to Equation (14), the required \( \mu_s \Omega_s \) value is of the order of that of a millisecond magnetar if \( \dot{q} \sim 10^{-5} \), achievable for a rapidly spinning BH. So the putative GBM signal associated with GW150914 could be interpreted with this model. There are suggestions that the GBM signal may not be real (e.g., Greiner et al. 2016; Xiong 2016). If so, one may place an upper limit on \( \dot{q} \) of the order of \( 10^{-5} \). The non-detection of \( \gamma \)-ray signals from LVT151012 and GW151226 (Racusin et al. 2011; Smartt et al. 2016) could pose an upper limit on \( \dot{q} \) to the same order.

The delay and the short duration of the GBM transient with respect to GW150914 could be readily explained. According to Equation (12), approximating \( M \sim 30 M_\odot \) for both BHs in GW150914, one may estimate \( \Delta t_{\text{d}} \lesssim 5 \) ms, which is \( \ll \) the delay timescale \( \Delta t_{\text{GRB}} \sim 0.4 \) s. One therefore has \( \Delta t_{\text{GRB}} \sim t_1 \) (noticing \( 1 + z \sim 1 \)), which gives a constraint on the onset radius of emission
\[ R_1 \sim 2\Gamma_2^2 c\tau_{\text{GRB}} = (2.4 \times 10^{14} \text{ cm}) \left( \frac{\Gamma_1}{100} \right)^2 \left( \frac{\Delta t_{\text{GRB}}}{0.4 \text{ s}} \right). \]  
(24)

The weak signal does not allow a precise measurement of \( t_r \) and \( t_d \). In any case, the pulse is asymmetric (Connaughton et al. 2016) with \( t_d = t_2 \gg t_1 = t_2 - t_1 \), consistent with the theory. The total duration is \( \tau = 2t_2 - t_1 \approx t_2 \), which defines the decay timescale due to the angular spreading curvature effect. One can then estimate the radius where emission ceases, i.e.,
\[ R_2 \sim 2\Gamma_2^2 c_2 t_2 \sim 2\Gamma_2^2 c_2 \tau = (6.0 \times 10^{14} \text{ cm}) \left( \frac{\Gamma_2}{100} \right)^2 \left( \frac{\tau}{1 \text{ s}} \right). \]  
(25)

Even though the Lorentz factor \( \Gamma \) for this kind of GRB is unknown, we can see that for nominal values (\( \Gamma_1 \sim \Gamma_2 \sim 100 \)) of known GRBs (Liang et al. 2010), the emission radius is much greater than the photosphere radius, suggesting that the GRB emission comes from an optically thin region. The large radius is consistent with the expectation of the models that invoke magnetic dissipation in a Poynting-flux-dominated outflow (Lyutikov & Blandford 2003; Zhang & Yan 2011).

6. EVENT RATE DENSITIES

For \( \dot{q} = 10^{-9} \)–\( 10^{-6} \) needed to produce FRBs, the required BH \( \mu_s \Omega_s \) is \( \sim (10^{23} \text{–} 10^{34}) \text{ G cm}^{-3} \text{ s}^{-1} \), which is much smaller than that of a millisecond magnetar. This suggests that a moderately spinning BH with a moderate magnetic field in a merger system could make an FRB. One would expect more associations of BH–BH mergers with FRBs than GRBs.

The inferred event rate density of BH–BH mergers from the detections of GW150914, GW151226, and LVT151012 (Abbott et al. 2016c) is \( \sim (9–240) \) Gpc\(^{-3} \text{ yr}^{-1} \). The FRB event rate density may be estimated as
\[ \rho_{\text{FRB}} = \frac{365N_{\text{FRB}}}{(4\pi/3)D_z^3} \approx (5.7 \times 10^3 \text{ Gpc}^{-3} \text{ yr}^{-1}) \times \left( \frac{D_z}{3.4 \text{ Gpc}} \right)^{3} \left( \frac{N_{\text{FRB}}}{2500} \right), \]  
(26)

where \( N_{\text{FRB}} \) is the daily all-sky FRB rate that is normalized to 2500 (Keane & Petroff 2015), and \( D_z \) is the comoving distance of the FRB normalized to 3.4 Gpc (\( z = 1 \)). One can see that the FRB rate is at least 20 times higher than the BH–BH merger rate (see also Callister et al. 2016). Recently Keane et al. (2016) claimed a cosmological origin of FRB 150418. Spitler et al. (2016), on the other hand, reported repeating bursts from FRB 121102, which point toward an origin of a young pulsar, probably in nearby galaxies (e.g., Connor et al. 2016; Cordes & Wasserman 2016). Based on radio survey data, Vedantham et al. (2016) suggested that the fraction of cosmological FRBs with bright radio afterglow as FRB 150418 should be a small fraction of the entire FRB population. Our analysis suggests that the BH–BH mergers can account for the cosmological FRBs if their fraction is less than 5% and if all BH–BH mergers can have \( \dot{q} \) at least \( 10^{-10} \)–\( 10^{-9} \). If the radio transient following FRB 150418 (Keane et al. 2016) is indeed the afterglow of the FRB (cf. Li & Zhang 2016; Williams & Berger 2016), then the observation is consistent with the prediction of this model (Zhang 2016).

7. SUMMARY AND DISCUSSION

For BH–BH mergers, if at least one of the BHs carries a certain amount of charge, the inspiral process generates a loop circuit, which induces a magnetic dipole. The rapid evolution of the magnetic moment of the system leads to a magnetospheric outflow with an increasing wind power. If \( \dot{q} \) can be as large as \( \sim (10^{-9} \text{–} 10^{-8}) \), the magnetospheric wind right before the coalescence may produce an FRB, and the BH–BH mergers may contribute to some cosmological FRBs. If \( \dot{q} \) could be as large as \( \sim (10^{-5} \text{–} 10^{-4}) \), a short-duration GRB may be produced. The putative short GRB signal associated with GW150914 (Connaughton et al. 2016) may be interpreted with this model.
The near-isotropic nature of the magnetosphere wind conjectured in this model suggests that every BH–BH merger should be accompanied by an EM counterpart (if $\hat{q}$ is large enough). The detection of an FRB (or even a GRB) associated with future BH–BH merger GW events would verify this model and lead to a measurement to $\hat{q}$ (since the luminosity is essentially a function of $\hat{q}$ only). The non-detections of GRBs and FRBs associated with these mergers, on the other hand, would place an upper limit on $\hat{q}$ allowed for astrophysical BHs.

The same physical picture naturally applies to NS–NS and NS–BH merger systems as well. Since those systems have at least one NS, it is guaranteed that at least one member of the merger system carries a $\hat{q}$ large enough to produce cosmological FRBs (see also Wang et al. 2016 for an alternative trigger mechanism). The detectable event rate of these mergers, however, is not much larger than BH–BH mergers, since in a large solid angle of such a merger, the FRB could not escape due the absorption of the dynamical ejecta launched during the merger. In systems with larger $\hat{q}$, the pre-merger dynamical magnetospheric activities would make a possible hard electromagnetic transient leading the main episode of the short GRB (see also a recent discussion on this aspect by Metzger & Zivancev 2016). A detection or an upper limit on this signal would give interesting constraints on the properties of the pre-merger systems.

I thank the referee, Anatoly Spitkovsky, for constructive comments and criticisms, and Mitch Begelman, Zi-Gao Dai, Tong Liu, Peter Mészáros, Kohta Murase, Martin Rees, Scott Tremaine, Z. Lucas Uhm, Shao-Lin Xiong, and Bin-Bin Zhang for helpful comments and discussion. This work is partially supported by NASA NNX15AK85G and NNX14AF85G.

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