The electromagnetic signals of compact binary mergers

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ABSTRACT
Compact binary mergers are prime sources of gravitational waves (GWs), targeted by current and next generation detectors. The question ‘what is the observable electromagnetic (EM) signature of a compact binary merger?’ is an intriguing one with crucial consequences to the quest for GWs. We present a large set of numerical simulations that focus on the EM signals that emerge from the dynamically ejected subrelativistic material. These outflows produce on a time-scale of a day macronovae – short-lived infrared (IR) to ultraviolet (UV) signals powered by radioactive decay. Like in regular supernovae the interaction of this outflow with the surrounding matter inevitably leads to a long-lasting remnant. We calculate the expected radio signals of these remnants on time-scales longer than a year, when the subrelativistic ejecta dominate the emission. We discuss their detectability in 1.4 GHz and 150 MHz and compare it with an updated estimate of the detectability of short gamma-ray bursts’ orphan afterglows (which are produced by a different component of this outflow). We find that mergers with characteristics similar to those of the Galactic neutron star binary population (similar masses and typical circummerger Galactic disc density of ∼1 cm−3) that take place at the detection horizon of advanced GW detectors (300 Mpc) yield 1.4 GHz [150 MHz] signals of ∼50 [300] μJy, for several years. The signal on time-scales of weeks is dominated by the mildly and/or ultrarelativistic outflow, which is not accounted for by our simulations, and is expected to be even brighter. Upcoming all sky surveys are expected to detect a few dozen, and possibly more, merger remnants at any given time thereby providing robust lower limits to the mergers rate even before the advanced GW detectors become operational. The macronovae signals from the same distance peak in the IR to UV range at an observed magnitude that may be as bright as 22–23 about 10 h after the merger but dimmer, redder and longer if the opacity is larger.

Key words: gravitational waves – stars: neutron – radio continuum: general – surveys – gamma-ray burst: general.

1 INTRODUCTION

Compact binary (neutron star–neutron star, ns2, or black hole–neutron star, nsh) mergers are prime sources of gravitational radiation. The gravitational wave (GW) detectors LIGO (Abbott et al. 2009a), Virgo (Acernese et al. 2008) and GEO600 (Grote 2008) are designed to optimally detect merger signals. These detectors have been operational intermittently during the last few years reaching their nominal design sensitivity (Abbott et al. 2009b; Sengupta 2010; The LIGO Scientific Collaboration & the Virgo Collaboration 2010) with detection horizons of a few dozen Mpc for ns2 and almost a hundred Mpc for nsh mergers (the LIGO–Virgo collaboration adopts an optimal canonical distance of 33/70 Mpc; Abadie et al. 2010). Both LIGO and Virgo are being upgraded now and by the end of 2015 are expected to be operational at sensitivities ∼10–15 times greater than the initial LIGO (Smith 2009), reaching a detection horizon of a few hundred Mpc for ns2 mergers and about a Gpc for nsh mergers (445/927 Mpc are adopted by the LIGO–Virgo collaboration as canonical values; Abadie et al. 2010).

Understanding the observable electromagnetic (EM) signature of compact binary mergers has several observational implications. First, once the detectors are operational it is likely that the first detection of a GW signal will be around or even below threshold. Detection of an accompanying EM signal will confirm the discovery, thereby increasing significantly the sensitivity of GW detectors.

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The EM signal that is often considered as the most promising counterpart to GWs is that of short gamma-ray bursts (GRBs), which are thought to arise from compact binary merger events (Eichler et al. 1989). The estimated rate of short GRBs is indeed comparable to binary pulsar estimates (Guetta & Piran 2006; Nakar, Gal-Yam & Fox 2006; Guetta & Stella 2009). However, while appealing, the association is not proven yet (Nakar 2007), and even if there is an association, short GRBs are observed only if their relativistic jets points towards us. If short GRBs are binary mergers then their observed rate, \( \sim 10 \text{Gpc}^{-3} \text{yr}^{-1} \), provides a lower limit to the merger rate. The true rate is higher and it depends on the poorly constrained beaming angle, which results in an uncertainty of almost two orders of magnitude. While a GRB that is observed off-axis is undetectable in \( \gamma \)-rays, it produces a long-lasting radio ‘orphan’ afterglow, which may be detectable (Rhoads 1997; Waxman, Kulkarni & Frail 1998; Frail, Waxman & Kulkarni 2000; Levinson et al. 2002; Gal-Yam et al. 2006). A key point in estimating the detectability of GRB orphan afterglows is that the well constrained observables are the isotropic equivalent energy of the flow and the rate of bursts that point towards earth. However, the detectability of the orphan afterglows depends only on the total energy and true rate, namely on the poorly constraint jet beaming angle. Levinson et al. (2002) have shown that while narrower beaming increases the true rate it reduces the total energy, and altogether reduces the detectability of radio orphan afterglows. This counterintuitive result makes the detectability of late emission from a decelerating jet, which produced a GRB when it was still ultrarelativistic, less promising.

Regardless of the amount of ultrarelativistic outflow that is launched by compact binary mergers, and of whether they produce short GRBs or not, mergers do launch energetic subrelativistic and mildly relativistic outflows (e.g. Rosswog et al. 1999; Ruffert & Janka 2001; Rosswog 2005; Rosswog & Price 2007; Yamamoto, Shibata & Taniguchi 2008; Rezzolla et al. 2010), unless the equation of state at supranuclear densities is extremely soft (Rosswog et al. 2000). Recently, Nakar & Piran (2011) showed that the interaction of these outflows with the surrounding matter will inevitably produce radio counterparts. An additional source of EM signal was synthesized, radioactive elements in the ejected debris from the merger can drive short-lived supernova (SN)-like event often referred to as ‘macronova’ (Kulkarni 2005). Metzger et al. (2010) find that if 0.01 Msun is ejected then the optical emission from a merger at 

\[ 300 \text{Mpc} \text{ peaks after } \sim 1 \text{d at } m_V \approx 23 \text{ mag. For a recent discussion of the detectability of the various EM counterparts of GW sources see Metzger & Berger (2012).} \]

In addition to EM signals, there will be a strong (\( \sim 10^{35} \text{ erg} \)) burst of \( \sim 10\text{ MeV} \) neutrinos, similar to what is produced by a core-collapse SN. But since compact binary mergers are orders of magnitude rarer than SN events, the chances of detecting neutrinos from a cosmological merger event are essentially zero.

In a companion paper (Rosswog, Nakar & Piran 2013, hereafter Paper I) we have investigated to which extent dynamical collisions as they occur, for example, in a globular cluster are different from a GW-driven compact binary merger. In this paper we concentrate entirely on binary mergers and we systematically explore the neutron star (ns) binary parameter space in a large set of simulations. We use numerical simulations of the merger process to find the properties of the dynamically ejected outflow for different masses of the coalescing compact stars. We then take the resulting ejecta profiles and calculate the EM transients (i.e. radio flares and macronovae) that are related to the dynamical ejecta of ns and nsh binaries.

Our study does not account for other types of outflows such as neutrino-driven winds (which yield moderate velocities of \( \sim 0.1c \); Dessart et al. 2009) or mildly and/or ultrarelativistic outflows that may emerge from close to the compact object at the centre of the merger. Since only the subrelativistic dynamically ejected material is explored we restrict the light-curve calculation to time-scales of a year and longer, when the subrelativistic component dominates the emission. On shorter time-scales of weeks and months the mildly relativistic component dominates. Since the radio luminosity depends strongly on the outflow velocity, emission on time-scales of weeks and months will be brighter than the one that we find here even if the mildly relativistic component carries a small fraction out of the total outflow energy (see Nakar & Piran 2011 for details). The radio flares depend sensitively on the surrounding interstellar medium (ISM) density. We focus here on physical parameters found in known Galactic ns binaries. Since all known binaries reside in the Galactic disc, we consider a uniform density of \( 1 \text{ cm}^{-3} \) (Draine 2011) as the most likely circummerger environment.

Our new simulations that focus on the ejecta also enable us to revise the estimates of the macronovae light curves. These light curves are determined mostly by three ingredients: (i) the total amount of the ejecta and their velocity structure that we calculate here; (ii) the energy input from radioactive decay for which we use the most accurate estimates to date by Korobkin et al. (2012) and (iii) the (poorly known) material opacity, taken from Metzger et al. (2010). At a finer level the light curve and the peak luminosity depend also on the velocity distribution as the emission from a given mass element moving with a specific velocity peaks when this element becomes optically thin. We use the simulation results to estimate this time-scale and in this way we obtain macronovae light curves for our different merger cases.

We begin (Section 2) with a brief discussion of our numerical simulations (more details can be found in Paper I). We focus in this discussion on the ejecta properties which are critical both for radio flares and for macronovae. In Section 3 we first provide an analytic calculation of the radio emission resulting from the interaction of single velocity ejecta with the surrounding ISM. Then, we expand this solution to an outflow with a distribution of velocities and we present a semi-analytic calculation to find the

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1 The current rate constraints on compact binary mergers are rather loose. The last LIGO and Virgo runs provided only weak upper limits on the merger rates: \( 1.4 \times 10^{4} \text{Myr}^{-1} (10^{10} \text{L}_\odot)^{-1} \) corresponding to \( \sim 2 \times 10^{5} \text{Gpc}^{-3} \text{yr}^{-1} \) for ns\(^2\) and \( 3600 \text{Myr}^{-1} (10^{10} \text{L}_\odot)^{-1} (\sim 5 \times 10^{4} \text{Gpc}^{-3} \text{yr}^{-1}) \) for nsh\(^2\) (Abbott et al. 2009b). Estimates based on the observed binary pulsars in the Galaxy are highly uncertain, with values in the range \( 20-2 \times 10^{4} \text{Gpc}^{-3} \text{yr}^{-1} \) (Narayan, Piran & Shemi 1991; Phinney 1991; Kalogera et al. 2004a,b; Abadie et al. 2010). There are no direct estimates of nshh mergers, as no such system has ever been observed, and here one has to rely only on a parameter-dependent population synthesis (e.g. Belczynski et al. 2008; Mandel & O’Shaughnessy 2010).

2 ns\(^2\) binaries are in random locations in the Galactic disc. Draine (2011) estimates that 0.4 of the disc volume has a density of about 0.6 cm\(^{-3}\), 0.1 of the volume has a much higher density while the rest has a low density.
Overview over the performed simulations. (see e.g. the data compilation in Lattimer & Prakash 2010).

\begin{tabular}{cccccccc}
Run & $m_1$ (M\odot) & $m_2$ (M\odot) & $N_{\text{SPH}}$ (10\textsuperscript{6}) & $t_{\text{end}}$ (ms) & $m_{\text{ej}}$ (M\odot) & $v_{\text{esc}}$ & $E_{\text{kin}}$ (10\textsuperscript{50} erg) \\
\hline
1 & 1.0 & 1.0 & 1.0 & 15.3 & $7.64 \times 10^{-3}$ & 0.10 & 1.0 \\
2 & 1.2 & 1.0 & 1.0 & 15.3 & $2.50 \times 10^{-2}$ & 0.11 & 3.4 \\
3 & 1.4 & 1.0 & 1.0 & 16.5 & $2.91 \times 10^{-2}$ & 0.13 & 4.8 \\
4 & 1.6 & 1.0 & 1.0 & 31.3 & $3.06 \times 10^{-2}$ & 0.13 & 5.3 \\
5 & 1.8 & 1.0 & 1.0 & 30.4\textsuperscript{a} & $>1.64 \times 10^{-2}$ & 0.13 & 3.2 \\
6 & 2.0 & 1.0 & 0.6 & 18.8\textsuperscript{a} & $>2.39 \times 10^{-2}$ & 0.16 & 6.0 \\
7 & 1.2 & 1.2 & 1.0 & 15.4 & $1.68 \times 10^{-2}$ & 0.11 & 2.3 \\
8 & 1.4 & 1.2 & 1.0 & 13.9 & $2.12 \times 10^{-2}$ & 0.12 & 3.2 \\
9 & 1.6 & 1.2 & 1.0 & 14.8 & $3.33 \times 10^{-2}$ & 0.13 & 6.2 \\
10 & 1.8 & 1.2 & 1.0 & 21.4 & $3.44 \times 10^{-2}$ & 0.14 & 7.0 \\
11 & 2.0 & 1.2 & 0.6 & 15.1\textsuperscript{a} & $>2.95 \times 10^{-2}$ & 0.14 & 6.0 \\
12 & 1.4 & 1.4 & 1.0 & 13.4 & $1.28 \times 10^{-2}$ & 0.10 & 1.6 \\
13 & 1.6 & 1.4 & 1.0 & 12.2 & $2.36 \times 10^{-2}$ & 0.12 & 4.0 \\
14 & 1.8 & 1.4 & 1.0 & 13.1 & $3.84 \times 10^{-2}$ & 0.14 & 7.6 \\
15 & 2.0 & 1.4 & 0.6 & 15.0 & $3.89 \times 10^{-2}$ & 0.15 & 8.7 \\
16 & 1.6 & 1.6 & 1.0 & 13.2 & $1.97 \times 10^{-2}$ & 0.11 & 2.9 \\
17 & 1.8 & 1.6 & 1.0 & 13.0 & $1.67 \times 10^{-2}$ & 0.12 & 2.7 \\
18 & 2.0 & 1.6 & 0.6 & 12.4 & $3.79 \times 10^{-2}$ & 0.14 & 7.6 \\
19 & 1.8 & 1.8 & 1.0 & 14.0 & $1.50 \times 10^{-2}$ & 0.12 & 2.8 \\
20 & 2.0 & 1.8 & 0.6 & 11.0 & $1.99 \times 10^{-2}$ & 0.13 & 3.7 \\
21 & 2.0 & 2.0 & 0.2 & 21.4 & $1.15 \times 10^{-2}$ & 0.11 & 1.8 \\
22 & 5\textsuperscript{a} & 1.4 & 0.2 & 138.7 & $2.38 \times 10^{-2}$ & 0.15 & 6.0 \\
23 & 10\textsuperscript{b} & 1.4 & 0.2 & 139.3 & $4.93 \times 10^{-2}$ & 0.18 & 18.2 \\
\end{tabular}

\textsuperscript{a}The secondary is still orbiting at the end of the computations.

\textsuperscript{b}The primary is a bh.

These findings are more than enough of a motivation for a broad scan of the ns binary parameter space. We explore ns masses between 1.0 and 2.0 M\odot in steps of 0.2 M\odot. All our ns have a negligible initial spin, consistent with the results of Bildsten & Cutler (1992) and Kochanek (1992). Our simulations make use of the smooth particle hydrodynamics (SPH) method, see Monaghan (2005) and Rosswog (2009) for recent reviews. Our code is an updated version of the one that was used in earlier studies (Rosswog & Davies 2002; Rosswog & Liebendörfer 2003; Rosswog, Ramirez-Ruiz & Davies 2003; Rosswog 2005). It uses the Shen et al. equation of state (EOS; Shen et al. 1998a,b), an opacity-dependent multi-flavour neutrino leakage scheme (Rosswog & Liebendörfer 2003) and a time-dependent artificial viscosity treatment, see Rosswog et al. (2000, 2008) for details. In all nsbh simulations Newtonian gravity was employed and the black hole (bh) was vested with an absorbing boundary at the Schwarzschild radius. All simulations used a simple GW emission back-reaction force (Davies et al. 1994). The performed simulations complement those that have been presented in Paper I. For completeness, we have also performed two simulations of nsbh binary systems. Since we found in earlier studies (Rosswog, Speith & Wynn 2004) cases of long-lived, episodic mass transfer, we started the two cases with lower numerical resolution that use the approximation of Newtonian gravity. Relativistic binary simulations with small mass ratios and large bh spin parameters may be particularly prone to a stable mass transfer phase, see Shibata & Taniguchi (2011) for a further discussion.

3 Phases of stable mass transfer are not necessarily restricted to simulations that use the approximation of Newtonian gravity. Relativistic binary simulations with small mass ratios and large bh spin parameters may be particularly prone to a stable mass transfer phase, see Shibata & Taniguchi (2011) for a further discussion.
In all investigated cases we find $\sim 10^{-2} \, M_\odot$ of unbound material ($E_{\text{kin}} + E_{\text{pot}} > 0$). A deviation of the mass ratio from unity has the tendency to enhance the amount of ejected mass, we find in contrast no clear tendency with the total system mass, see column six in Table 1. In all cases the ejecta velocities are below $0.5c$, their mass-averaged values are given in column seven, they are typically close to $0.13c$.

Figs 1 and 2 depict the remnant structures in the orbital (XY) plane and perpendicular to it (XZ) at the end of the simulations. Clearly, visible is the sensitivity to deviations from a mass ratio of...
unity: even differences of 15 per cent in the stellar masses lead to large asymmetries, i.e. one pronounced tidal tail rather than a disc resulting from initially two such tails. The debris matter cannot cool rapidly enough, therefore, it is puffed up, see Fig. 2, but at the end of the simulations the configuration is far from spherical symmetry. Although this may lead to some viewing angle dependencies, we assume in the following for simplicity spherically symmetric outflows.

Compact binary mergers inevitably eject mass in various forms.

(a) If compact binaries indeed power short GRBs, they have to launch ultrarelativistic outflows with Lorentz factors $>100$–$1000$. There are several mechanisms by which this could be achieved (e.g. Blandford & Znajek 1977; Piran 2004; Hawley & Krolik 2006; Lee & Ramirez-Ruiz 2007; Nakar 2007).

(b) Once the merger has happened, the hot remnant emits a few times $10^{53}$ erg s$^{-1}$ in $\sim 20$ MeV neutrinos (Eichler et al. 1989; Ruffert & Janka 2001; Rosswog & Liebendörfer 2003; Sekiguchi 2011). The peak of the neutrino emission is delayed with respect to the merger itself by the time it takes to form a hot accretion torus, about 10 ms. At these huge luminosities the neutrinos drive a strong baryonic wind of $M \sim 10^{-3} M_\odot$ s$^{-1}$ and $v \sim 0.1c$ (Dessart et al. 2009), preferentially in the direction of the rotation axis.

(c) It is very likely that the outflows of process (a) and (b) interact near the rotation axis/the inner disc. Such an interaction plausibly produces moderately relativistic matter outflows near the jet–wind interface.

(d) As the accretion disc evolves it spreads viscously until dissipational heating and/or the recombination of nucleons into nuclei unbind a large fraction of the late-time disc (Chen & Beloborodov 2007; Lee & Ramirez-Ruiz 2007; Beloborodov 2008; Metzger, Piro & Quataert 2008).

(e) Gravitational torques dynamically eject matter directly at first contact with velocities $>0.1c$, see Table 1.

While all of the above mass loss processes undoubtedly occur, the quantitative calculation of processes (a)–(d) is technically very demanding and the relevant physical processes are not included in the presented simulations. We therefore do not consider here their contribution to the EM signature (which will be most important for the radio emission at early times – as discussed later). Instead, we focus entirely on the signature of the dynamic ejecta, which can be reliably calculated.\footnote{Note that the presented ejecta amounts are numerically fully converged. They can, however, depend on the included physics ingredients such as the EOS or the treatment of gravity [Newtonian versus general relativity (GR)].} Their properties are entirely set during the first contact and by the end of the simulation the unbound material is moving ballistically, with fast moving ejecta ahead of slower one. The ballistic motion ends only when the dynamical ejecta start being decelerated by the ambient material months–years after the merger (see below). The dynamically ejected debris is expected to be the most energetic of all the outflow components. As a result its velocity profile is not expected to be significantly affected by interaction with other outflow components and its profile at the end of the numerical simulation provides a good approximation of the initial conditions for the EM signal calculations. Fig. 3 depicts the dynamically ejected mass above a given velocity for the different merger cases that we considered. Typical values range from 0.01 $M_\odot$ to 0.05 $M_\odot$ with $v > 0.1-0.16c$.

\[ v_{\text{dec}} = \frac{3E}{4\pi m c^2} \beta_0^{1/3} \approx 10^{17} \text{ cm s}^{-1}, \frac{n^{-1/3}}{\beta_0^{2/3}} \eta_{17} \]

\[ t_{\text{dec}} = \frac{R_{\text{dec}}}{c\beta_0} \approx 30 d \frac{E_{49}^{1/3} n^{-1/3}}{\beta_0^{5/3}}, \]

where we approximate $\Gamma_0 - 1 \approx \beta_0^2$ and ignore relativistic effects. Here and in the following, unless stated otherwise, $\eta$ denotes the value of $\eta/10^4$ in cgs units. At a radius $R > R_{\text{dec}}$ the flow decelerates assuming the Sedov–Taylor self-similar solution, so the outflow velocity can be approximated as

\[ \beta = \beta_0 \left( \frac{1}{(R/R_{\text{dec}})} \right)^{3/2} \frac{R \leq R_{\text{dec}}}{R > R_{\text{dec}}}. \]

If the outflow is collimated, highly relativistic and points away from a generic observer, as will typically happen if the mergers produce short GRBs, the emission during the relativistic phase will be suppressed by relativistic beaming. Observable emission is produced only once the external shock decelerates to mildly relativistic...
velocities and the blast wave becomes quasi-spherical. This takes place when \( \Gamma \approx 2 \) namely at \( \tau_{\text{dec}}(\beta_0 = 1) \). From this radius the hydrodynamics and the radiation become comparable to that of a spherical outflow with an initial Lorentz factor \( \Gamma_0 \approx 2 \). This behaviour is the source of the late radio GRB orphan afterglows (Rhoads 1997; Levinson et al. 2002). Our theory is therefore applicable for the detectability of mildly and non-relativistic outflows as well as for radio orphan GRB afterglows.

Emission from Newtonian and mildly relativistic shocks is observed in radio SNe and late phases of GRB afterglows. These observations are well explained by a theoretical model involving synchrotron emission of shock accelerated electrons in an amplified magnetic field. The success of this model in explaining the detailed observations of radio Ib/c SNe (e.g. Chevalier 1998; Soderberg et al. 2005; Chevalier & Fransson 2006) allows us to employ the same microphysics here. Energy considerations show that both the electrons and the magnetic field carry significant fractions of the total internal energy of the shocked gas, \( \epsilon_e \approx \epsilon_B \approx 0.1 \). These values are consistent with those inferred from late radio afterglows of long GRBs (e.g. Frail et al. 2000, 2005). The observed spectra indicate that the distribution of the accelerated electrons’ Lorentz factors, \( \gamma \), is a power law \( dN/d\gamma \propto \gamma^{-x} \) at \( \gamma > \gamma_\text{m} \), where \( p \approx 2.1-2.5 \) in mildly relativistic shocks (e.g. the radio emission from GRB associated SNe and late GRB afterglows) and \( p \approx 2.5-3.0 \) in Newtonian shocks (as seen in typical radio SNe; Chevalier 1998, and references therein). The value of \( \gamma_\text{m} \) is not observed directly but it can be calculated based on the total energy of the accelerated electrons, \( \gamma_\text{m} \approx \beta_0^{-2} \frac{m_{\text{e}}^2 c^2}{m_\text{B}} \).

The radio spectrum generated by the shock is determined by two characteristic frequencies.\(^3\) One is

\[
v_\text{m} \approx 1 \text{ GHz} n^{1/2} \beta_0^{-1} \epsilon_{\text{B},-1}^{-2/3} \epsilon_{e,-1}^{2/3},
\]

(4)

the typical synchrotron frequency of electrons with the typical (also minimal) Lorentz factor \( \gamma_\text{m} \). The other is \( v_\text{obs} \), the synchrotron self-absorption frequency. We show below that since we are interested in the maximal flux at a given observed frequency, \( v_\text{obs} \) may play a role only if it is larger than \( v_\text{m} \). Its value in that case is

\[
v_\text{obs} (> v_\text{m}) \approx 1 \text{ GHz} R_{17}^{1/7} n^{6/7} \beta_0^{4/7} \epsilon_{\text{B},-1}^{2/7} \epsilon_{e,-1}^{1/7} \beta^{1/7}.
\]

(5)

Consider now a given observed frequency \( v_{\text{obs}} \). We are interested in the light curve near the peak flux at this frequency. There are three possible types of light curves near the peak corresponding to (i) \( v_{\text{obs}} > v_{\text{m}} \), (ii) \( v_{\text{obs}} < v_{\text{m}} < v_{\text{obs}} \), and (iii) \( v_{\text{obs}} < v_{\text{eq}} \). Where we define

\[
v_{\text{eq}} = 1 \text{ GHz} E_{49}^{4/7} n^{4/7} \beta_0^{-1/7} \epsilon_{\text{B},-1}^{1/7} \epsilon_{e,-1}^{1/7},
\]

(10)

as the frequency at which \( \gamma_{\text{eq}} = \gamma_{\text{m}} \). In Fig. 5 we show a sketch of the time evolution of \( v_{\text{a}} \) and \( v_{\text{m}} \) and the corresponding ranges of \( v_{\text{obs}} \) in which each of the cases is observed.

To estimate the time and value of the peak flux we recall, that at all frequencies the flux increases until \( \tau_{\text{dec}} \). In case (i), \( v_{\text{obs}} > v_{\text{m}} \), the deceleration time, \( \tau_{\text{dec}} \) is also the time of the peak.

The reason is that while \( F_{\text{m}} \) increases, \( v_{\text{obs}} \) decreases fast enough so that \( F_{\text{obs}} \) decreases after \( \tau_{\text{dec}} \). Note that in that case \( v_{\text{a}} \) plays no role since it decreases after deceleration. Overall, in this case the flux peaks at \( \tau_{\text{dec}} \) and \( F_{\text{obs}, \text{peak}} = F_{\text{dec}}(v_{\text{obs}}/v_{\text{dec}})^{-1(p-1)/2} \).

In the two other cases, (ii) and (iii), \( v_{\text{obs}} \approx v_{\text{m}} \), and/or \( v_{\text{obs}} \approx v_{\text{eq}} \), the flux keeps rising at \( t > \tau_{\text{dec}} \) until \( v_{\text{obs}} = v_{\text{eq}}(t) \) whichever comes last. To find out which one of the two frequencies is it, we compare \( v_{\text{obs}} \) with \( v_{\text{eq}} \). At \( t > \tau_{\text{dec}} \), \( v_{\text{obs}} \) decreases faster than \( v_{\text{eq}} \). Therefore in case (ii) where \( v_{\text{eq}} < v_{\text{obs}} \), the last frequency to cross \( v_{\text{obs}} \) is \( v_{\text{a}} \), and the peak flux is observed when \( v_{\text{obs}} = v_{\text{a}}(t) \). In case (iii) where \( v_{\text{obs}} < v_{\text{eq}} \), the last frequency to cross \( v_{\text{obs}} \) is \( v_{\text{obs}} \), and the peak flux is observed when \( v_{\text{obs}} = v_{\text{obs}}(t) \). Now, it is straightforward to calculate the peak flux, \( F_{\text{obs, peak}} \) and the time that it is observed, \( \tau_{\text{peak}} \). For different frequencies. It is also straightforward to calculate the flux temporal evolution prior and after \( \tau_{\text{dec}} \) using equations (4)–(6) and the relation \( t \propto R \) which holds at \( t < \tau_{\text{dec}} \) and \( \beta \propto t^{-3/5} \) at \( t > \tau_{\text{dec}} \). The peak fluxes, the times

\[6\] Note that if \( v_{\text{dec}} < v_{\text{a}} \), this equality will never take place. In that case \( v_{\text{eq}} \) is the frequency at which this equality would have happened if \( \Gamma_0 \) would have been large enough (see Fig. 5).
of the peak and the temporal evolution of the three different cases are summarized in Table 2. The overall different light curves are depicted in Fig. 6.

The most sensitive radio facilities are at frequencies of 1.4 GHz and higher. Equations (7) and (8) imply that in this frequency range, for most realistic scenarios, it is a case (i) light curve, i.e. \( v_a, \delta v_m < v_{\text{obs}} \). Therefore, Newtonian and mildly relativistic outflows as well as relativistic GRB orphan afterglows peak at \( t_{\text{dec}} \) with (Nakar & Piran 2011)

\[
F_{\text{vobs, peak}}(v_{\text{obs}} < v_{\text{eq}}, v_a, \delta v_m < v_{\text{obs}}) \approx 0.3 \text{ mJy } E_{\text{GRB}} \left( \frac{v_{\text{obs}}}{1.4 \text{ GHz}} \right)^{-\frac{n+1}{2}}.
\]

The regime of \( F_{\text{vobs, peak}} \) at lower radio frequencies (<1 GHz) depends on the various parameters. If the outflow is Newtonian or the density is low or the energy is low then \( v_a, \delta v_m < v_{\text{obs}}, v_{\text{dec}} < 100 \text{ MHz} \) and equation (11) is applicable. Otherwise low radio frequencies are in regime (iii), i.e. \( v_{\text{obs}} < v_{\text{eq}}, v_a, \delta v_m \). The flux peaks in this case at

\[
t_{\text{peak}}(v_{\text{obs}} < v_{\text{eq}}, v_a, \delta v_m) \approx 200 d E_{\text{GRB}} n \frac{5}{3} \varepsilon_{B, -1} \varepsilon_{e, -1}^{-1} \left( \frac{v_{\text{obs}}}{150 \text{ MHz}} \right)^{12} \nu_{\text{peak}}^{-2}.
\]

Figure 5. A sketch of the time evolution of \( v_a \) and \( v_m \) in two cases, \( v_a, \delta v_m < v_{\text{dec}} \) (top) and \( v_a, \delta v_m > v_{\text{dec}} \) (bottom). Also marked is the value of \( v_{\text{eq}} \). The vertical dashed line marks \( t_{\text{dec}} \). The ranges of \( v_{\text{obs}} \) at each of the cases is observed is separated by horizontal dashed lines and marked on the right. Note that in the bottom panel \( v_m \) and \( v_a \) are not crossing each other at \( t > t_{\text{dec}} \) and only two types of light curves, cases (i) and (iii), can be observed.

Figure 6. Schematic light curves of the three cases. The rising phase, marked in dashed line for each of the phases, is that of the last temporal power-law segment before the peak. After the peak all cases show the same power-law decay.

with

\[
F_{\text{vobs, peak}}(v_{\text{obs}} < v_{\text{eq}}, v_a, \delta v_m) \approx 50 \mu \text{Jy } E_{\text{GRB}}^{5/2} \varepsilon_{B, -1}^{5/2} \varepsilon_{e, -1}^{-1} \nu_{\text{obs}}^{13/2} \left( \frac{v_{\text{obs}}}{150 \text{ MHz}} \right)^{12}.
\]

In the last two equations we used \( p = 2.5 \) (other \( p \) values in the range 2.1–3 yield slightly different numerical factors and power laws).

To date, the best observed signal from a mildly relativistic blast wave is the radio emission that follows GRB associated SNe. The main difference is that in these cases the circumburst medium is typically a wind (i.e. \( n \propto R^{-2} \)) and therefore the density at early times is much larger than in the ISM and self-absorption plays the main role in determining the light curve. A good example for comparison of equation (11) with observations is the light curve of SN 1998bw. This light curve is observed at several frequencies at many epochs, enabling a detailed modelling that results in tight constraints of the blast wave and microphysical parameters. Li & Chevalier (1999) find that at the time of the peak at 1.4 GHz, about 40 d after the SN, taking \( \varepsilon_e = \varepsilon_B = 0.1 \), the energy in the blast wave is \( \sim 10^{50} \text{ erg} \), its Lorentz factor is \( \sim 2 \) and the external density at the shock radius is \( n \sim 1 \text{ cm}^{-3} \). The peak is observed when \( v_{\text{m}, \delta v_{\text{m}}} \leq v_{\text{obs}} \) and it depends only on these parameters (it is only weakly sensitive to the density profile). Therefore, equation (11) is applicable in that case. Indeed, plugging these numbers into equation (11) we obtain a flux of 20 mJy at the distance of SN 1998bw (40 Mpc), compared to the observed flux of 30 mJy.

Table 2. The observed flux before and after \( t_{\text{peak}} \) in the three different regimes.

| Case | \( F_{\text{vobs, peak}}/F_{\text{m, dec}} \) | \( t_{\text{peak}}/t_{\text{dec}} \) | \( F_{\text{vobs, peak}}/F_{\text{m, dec}} \) | \( t < t_{\text{peak}} \) |
|------|---------------------------------|----------------|---------------------------------|----------------|
| (i)  | \( v_{\text{m}, \delta v_{\text{m}}}, v_a < v_{\text{obs}} \) | \( (v_{\text{obs}}/v_{\text{m}, \delta v_{\text{m}}})^{-5/2} \) | 1 | \( \propto \nu_{\text{obs}}^{13/2} \) |
| (ii) | \( v_{\text{eq}} < v_{\text{obs}} < v_{\text{m}, \delta v_{\text{m}}} \) | \( (v_{\text{obs}}/v_{\text{m}, \delta v_{\text{m}}})^{-1/3} \) | \( (v_{\text{obs}}/v_{\text{m}, \delta v_{\text{m}}})^{-1/3} \) | \( \propto \nu_{\text{obs}}^{-2} \) |
| (iii) | \( v_{\text{obs}} < v_{\text{eq}}, v_a, \delta v_m \) | \( \frac{1}{2} \frac{1}{\nu_{\text{req}}/v_{\text{m}, \delta v_{\text{m}}} - 2} \frac{1}{\nu_{\text{req}}/v_{\text{m}, \delta v_{\text{m}}} - 2} \frac{1}{\nu_{\text{obs}}/v_{\text{m}, \delta v_{\text{m}}} - 2} \frac{1}{\nu_{\text{obs}}/v_{\text{m}, \delta v_{\text{m}}} - 2} \) | \( \propto \nu_{\text{obs}}^{13/2} \) | \( \propto \nu_{\text{obs}}^{-2} \) |

*The temporal evolution only during the last power-law segment before \( t_{\text{peak}} \). At earlier times the temporal evolution may be different.

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This is not surprising given that the model we use is based on that of radio SNe.

3.2 An outflow with a power-law velocity profile

The numerical simulations provide profiles of the outflow energy as a function of the velocity. Similarly to the case of an outflow with a single velocity we approximate the blast wave to be spherical. The main difference, when a range of velocities is considered, is that in addition to the forward shock, which is driven into the circumburst medium, there is a continuous reverse shock that is driven into the ejecta. If most of the outflow energy resides in low velocities then this reverse shock drives an increasing amount of energy into the shocked region, mitigating the deceleration of the forward shock. In that case the mass collected by the forward shock, $M(R)$, when its velocity is $\beta$ and its radius is $R$, is comparable to the mass in the ejecta with velocity $\geq \beta$. Thus the relation between the radius and the velocity can be found at any time by equating

$$M(R)(\beta c)^2 = E(\beta),$$

where this equality is correct up to a factor of order unity that depends on the exact ejecta profile. Chevalier (1982) calculates self-similar solutions for an outflow with a power-law distribution $\rho \propto \beta^{-k}$ in which most of the energy is at low velocities, i.e. for $k > 5$ since $E(>\beta) \propto \beta^{-(k-5)}$. Chevalier (1982) provides the exact coefficients of the solution for several values of $k$, all providing corrections of order unity to equation (14).

In Section 3.3 we apply equation (14) to the results of the numerical simulations and we calculate the resulting light curves from the ejecta profiles. However before doing so, we use equation (14) to find $R(t)$ and $\beta(t)$ for an outflow with a power-law velocity profile and we find an analytic solution for the flux for a power-law velocity profile. Consider an outflow with a power-law velocity profile and a minimal velocity $\beta_{\text{min}}$ and a total energy $E(>\beta_{\text{min}}) = E_{\text{tot}}$ that propagates in a constant density medium. Equation (14) implies that its radius and velocity evolve with time as

$$R = \left(\frac{3k^{-3}(\beta_{\text{min}} c)^{-3}E_{\text{tot}}}{4\pi(k-3)^{3/2}n_\rho}ight)^{1/4} t^{1/2},$$

$$\beta = \left(\frac{3(k-3)^3(\beta_{\text{min}} c)^{-3}E_{\text{tot}}}{4\pi k^{-3}n_\rho}\right)^{1/4} t^{-1/4},$$

plugging these values into equations (4)–(6) we find that for relevant parameters $v_e$, $v_m < 1.4$ GHz and

$$F_e(v > v_e, v_m) \approx 50 \mu Jy \exp^{-3.44(p-2.5)} E_{\text{tot},50} \frac{\beta_{\text{min}}}{0.2} \frac{n^{1.1}}{\epsilon_\gamma^{-1/2} t_{\text{year}}^{-5/2}} \frac{d_{5}^{2.5}}{d_{27}^{2} t_{\text{year}}^{-2} \epsilon_{\gamma}^{-3/2}},$$

$$= 50 \mu Jy \exp^{0.86 \left(\frac{\beta_{\text{min}}}{0.2}\right)^{3.44} n^{1.1} \epsilon_\gamma^{-1} t_{\text{year}}^{-5/2} d_{27}^{2} t_{\text{year}}^{-12} \epsilon_{\gamma}^{-3/2}} (k = 9; \ p = 2.5).$$

This equation is applicable starting at $t_{\text{dec}}$ that corresponds to the maximal velocity, $\beta_{\text{max}}$, and the energy that is carried by the material that moves at the maximal velocity (i.e. $E_{\beta_{\text{max}}} \approx E_{\text{tot}}(\beta_{\text{max}}/\beta_{\text{min}})^{-k-5}$) and up to $t_{\text{dec}}$ that corresponds to $\beta_{\text{min}}$ and $E_{\text{min}}$. At earlier times then $t_{\text{dec}}(\beta_{\text{max}} E_{\text{tot}})$ the light curve is the one described in Table 2 (with $\beta_{\text{max}} E_{\text{tot}}$) and at later times then $t_{\text{dec}}(\beta_{\text{min}} E_{\text{tot}})$ it joins the decaying light curve described in Table 2 ($F_e \propto t^{21-15\phi_{10}/10}$).

3.3 Radio light curves from the simulated outflows

To calculate the EM signatures we use the ejecta velocity profiles from the simulations and apply the approximations of Section 3.2. We use equation (14) to find $R(t)$ and $\beta(t)$ and subsequently plug them into equations (4)–(6) to calculate the light curve. We approximate the ejecta–ambient medium interaction as a spherical blast wave that propagates into a constant density, $n$. Behind the shock constant fractions of the internal energy, $\epsilon_\gamma = 0.1$ and $\epsilon_\gamma = 0.1$, are deposited in relativistic electrons and in magnetic field, where the electrons are accelerated to a power law with an index $p = 2.5$.

Inspection of the various ns–5 M$\odot$ mergers simulations reveal relatively large kinetic energies (at least $10^{50}$ erg and at times near $10^{51}$ erg, see Table 1) but with relatively low average velocities, around $0.1c$–$0.16c$. The corresponding peak fluxes and durations correspond, therefore, to the subrelativistic case examined in Nakar & Piran (2011). We calculate, for these simulated mergers, the light curves at two frequencies, 1.4 GHz and 150 MHz, and for two values of external densities $n = 1$ and $0.1$ cm$^{-3}$. The flux normalization is for events at a distance of $10^{27}$ cm, roughly at the detection horizon for ns$^2$ mergers by advanced LIGO and Virgo.

Here we calculate the radio emission only from the dynamical component of the ejecta (see Section 2). This component is launched preferentially along the equatorial plane, while faster moving outflows (related to processes ‘a’ and ‘c’ discussed in Section 2) are launched along the rotation axis. As a result the fast moving ejecta can propagate to large distances, ahead of the dynamical ejecta, and interact with the ambient medium. The faster moving components are expected to carry less energy than the slow moving ejecta. Yet, the strong dependence of the radio peak flux on the outflow velocity (equation 11) and the short deceleration time of the fast moving ejecta (weeks to months) imply that it is expected to dominate radio emission before the dynamical ejecta start decelerating (months to years). Thus, our estimates in this section are only lower limits on the true radio emission at early time (up to about a year for $n = 1$ cm$^{-3}$ and about 3 yr for $n = 0.1$ cm$^{-3}$). A glimpse of the influence of a mildly relativistic material can be seen in Figs 11 and 12 below which depict the light curves of nsbh mergers. Our simulations find that a mildly relativistic material is ejected in the 1.4 M$\odot$ ns–5 M$\odot$ bh case. Although it carries a much lower energy than the slower moving ejecta its strong effect is seen as early times. Since we expect the dynamical ejecta to be the most massive and most energetic of all the outflow components, it is not expected to be affected significantly by possible interaction between the different outflow components. If this is the case then our radio predictions are good approximations to the true emission at earlier times. If, however, the fast outflow is more energetic than the slow one, then also our late radio predictions are only lower limits of the true emission.

\footnote{We assume that the maximal velocity ejecta is still mildly relativistic, i.e. $\gamma_{\text{max}} \beta_{\text{max}} \lesssim 1$.}
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Figure 7. Radio light curves generated by interaction of the dynamically ejected subrelativistic outflows from equal mass ns\(^2\) mergers at 150 MHz and 1.4 GHz for 1 cm\(^{-3}\) circummerger density, \(\epsilon_B = \epsilon_e = 0.1\) and \(p = 2.5\). The merger distance is \(10^{27}\) cm, roughly the detection horizon for ns\(^2\) mergers by advanced LIGO and Virgo. The shaded region at \(t < 1\) yr reflect the fact that at this time the emission is expected to be dominated by mildly relativistic outflows, which are not included in our simulations.

The resulting light curves are shown in Figs 7–12 and the basic properties of the radio flares are summarized in Table 3. Figs 7 and 8 depict the radio light curves of the various ns\(^2\) mergers for \(n = 1\) cm\(^{-3}\). Our canonical case, a merger of two 1.4 M\(_\odot\) ns is marked in both figures with a thick solid line. The flux of the canonical merger is 0.04 mJy (0.2 mJy) at 1.4 GHz (150 MHz) 1–4 yr after the merger. It shows almost minimal emission among all equal mass mergers (Fig. 7), the emission increases by a factor of 5 (at 1.4 GHz) for a 1.8–1.8 M\(_\odot\) pair and it is almost similar for 1.0–1.0 M\(_\odot\). Fig. 8 depicts mergers with combinations of different ns masses. In mergers with a larger mass difference the secondary is disrupted more completely, producing a more prominent tidal tail. Therefore, the ejected mass rises with the mass ratio (see Table 1 and Fig. 1). As a result, the dynamically ejected outflows from unequal mass mergers, produce in general brighter and longer lived radio remnants. Indeed, in most cases of unequal mass mergers the luminosities, both in the 1.4 GHz and at 150 MHz, are brighter at late times than the equal mass case. In some cases they reach 0.3 mJy (2 mJy) a few years after the merger. Note that even a small mass difference of less than 15 per cent can make a large difference in the observed luminosity. For example, a 1.4–1.2 M\(_\odot\) merger produces twice as bright and twice as long remnant compared to a 1.4–1.4 M\(_\odot\) merger.

The external density is a critical parameter. Figs 9 and 10 depict the resulting light curves when the density is \(n = 0.1\) cm\(^{-3}\) for \(t > 3\) yr. In that case the signal, at the two considered frequencies, is almost an order of magnitude fainter in all cases, compared with \(n = 1\) cm\(^{-3}\). At 1.4 GHz the signal also evolves slower. Since the self-absorption frequency is below 150 MHz in that case, the light curves at both 1.4 GHz and 150 MHz have quite similar shapes.

Mergers of a bh and a ns eject even larger amounts of mass. Their signals, shown in Figs 11 and 12, are stronger reaching 1 mJy.
4.1 Detectability

We use the radio light curves of our canonical nst merger (1.4–1.4 M⊙) to estimate their detectability by current and future radio facilities. The estimates based on this merger case are conservative. First, since other merger cases, and especially those with even a minor mass difference between the coalescing stars, are brighter and second, since our simulations do not include all outflow sources.

We also consider here the detectability of nst merger based on the light curves calculated for a 1.4–10 M⊙ merger. Table 4 shows the detection horizons for these two merger scenarios for different radio telescopes and two external densities (n = 1 and 0.1 cm⁻³).

The radio flares depend sensitively on the surrounding circummerger density. Current robust knowledge concerning nst binaries arises from the observed Galactic population. These binaries are all observed in the Galactic disc whose typical density is about 1 cm⁻³ (Draine 2011). Hence we expect that this is the most relevant density to consider. In particular we stress that the estimated rate of nst that are based on these binaries (and used later when we estimate the detection rate of these radio flares) are all relevant for this particular population. An additional hypothetical population of nst mergers that take place in a much lower density environment and might produce weaker signals is not included in our radio counterpart estimates.

A nst merger that takes place in n = 1 cm⁻³ environment produces a radio remnant that can be easily observed by Extended Very Large Array (EVLA; 1 h integration) all the way to the advanced LIGO/Virgo detection horizon, with a rise time of ~1 yr and a decay time of several years. However, given the relative narrow EVLA field of view it requires either a ~10⁶ deg⁻² localization or a dedicated search in a larger error box. Given that the planned localization of advanced GW detectors is ~10–100 deg⁻² for events with a detector network signal-to-noise ratio of 10 (Wen & Chen 2010), a targeted search in the error box of a GW signal is certainly feasible. For that purpose other facilities with lower sensitivity but significantly larger field-of-view [e.g. Australian Square Kilometre Array Pathfinder (ASKAP)] may be more appropriate.

A nst merger in n = 0.1 cm⁻³ environment requires a more dedicated effort (10 h) in order to detect the resulting signal out to 300 Mpc by the EVLA. The rise and decay time are significantly longer (years and a decade, respectively). This makes the search for nst transients, even in the case of a ~10–100 deg⁻² GW localization, very challenging, unless they produce at least 10⁹⁹ erg in mildly relativistic outflow (which is not accounted for by our simulations).

nst mergers produce significantly brighter signals, and are therefore detectable up to a greater distances, ~1 Gpc, by the EVLA and near future facilities, even for density of n = 0.1 cm⁻³. This distance is similar to the detection horizon of these events by the advanced GW detectors. This makes nst mergers a promising target, however, the observed rate depends of course on their unknown merger rate.

Next we consider detectability in a blind radio survey. The number of radio remnants in a single snapshot all sky radio image is 

\[ N_{\text{all-sky}} = \mathcal{R} V \Delta t, \]

where \( V \) is the detectable volume, \( \Delta t \) is the time that the flux is above the detection limit and \( \mathcal{R} \) is the event rate. Assuming a single velocity outflow with typical parameters the 1.4 GHz light curve is a case (i) (see Section 3). Thus, we use equations (2) and (11), and the approximation \( \Delta t \approx \tau_{\text{rec}} \), to find that the number of radio coalescence remnants in a single 1.4 GHz whole sky snapshot with a detection limit \( \mathcal{F} \) (see also supplementary material in Nakar & Piran 2011):

\[ N_{\text{all-sky}}^{1.4 \text{GHz}} \approx 20 E_{49}^{11/6} n^{3/2} \left( \frac{h \nu}{\text{GHz}} \right)^{3/2} \frac{d}{\text{yr}^{-1}} \frac{I_{\text{rec}}}{\text{Jy cm}^{-2} \text{Hz}^{-1}} \left( \frac{\rho_0}{10^{-11} \text{g cm}^{-3}} \right) \mathcal{R}_{300} \mathcal{F}^{-3/2}, \]

where \( \mathcal{F} \) is the flux limit/0.1 mJy and \( \mathcal{R}_{300} \) is the merger rate in units of 300 Gpc⁻¹ yr⁻¹. Plugging values from Table 1 for the different runs we can obtain a first estimate, including the dependence on the circummerger density and microphysical parameters, of the number of all sky flares that are detectable. A better estimate for specific parameters can be obtained using the light curves we calculated in.
Section 3.3. Again we take 1.4–1.4 M⊙ as the canonical ns² merger and 1.4–10 M⊙ as the canonical nsbh merger.

A 1.4–1.4 M⊙ ns² merger that takes place in $n = 1$ cm$^{-3}$ environment is brighter than 0.1 mJy at 1.4 GHz for about 4 yr at a distance of 150 Mpc. Therefore, if as suggested by Galactic ns², this is the density of a typical merger environment, then $N_{\text{all-sky}}^{\text{ns}} \sim 20R_{\text{nsbh}}^3 R^{-3/2}$, if however, typical ns² merger takes place at $n = 0.1$ cm$^{-3}$ environment the number of remnants in an all sky snapshot drops by an order of magnitude.

A 1.4–10 M⊙ nsbh merger that takes place in $n = 1$ cm$^{-3}$ environment is brighter than 0.1 mJy at 1.4 GHz for about 4 yr at a distance of 1 Gpc, implying $N_{\text{all-sky}}^{\text{nsbh}} \sim 20R_{\text{nsbh}}^3 R^{-3/2}$, where $R_{\text{nsbh}}$ is the nsbh merger rate in units of Gpc$^{-3}$ yr$^{-1}$. Again, if the typical nsbh circummerger density is $n = 0.1$ cm$^{-3}$ then $N_{\text{all-sky}}$ drops by an order of magnitude.

To conclude, our calculations here strengthen the results of Nakar & Piran (2011) that there is a fair chance to detect merger radio remnants in a sub-mJy survey of even a part of the sky and a high chance to detect them in a whole sky survey. Such survey’s of a small portion of the sky are planned already with the EVLA, while a sub-mJy large-scale transient survey is part of the ASKAP Survey Science Projects. Finally, a 1 mJy 150 MHz survey with Low-Frequency Array (LOFAR) will find a comparable number of remnants, but these will have a longer rise time, and therefore will be harder to identify.

4.2 A comparison with short GRBs’ radio orphan afterglows

The outflows of short GRBs begin highly relativistic and probably highly beamed. Eventually they slow down and become detectable from all directions (see Section 3). Therefore, the rate estimate equation (18) is also applicable for radio orphan afterglows when $\beta_0 = 1$. However some of the parameters in equation (18) are not directly observable. The observed quantities are the isotropic equivalent $\gamma$-ray energy, $E_{\gamma,\text{iso}}$, and the rate of bursts that point to the observer $R_{\text{obs}}^{\text{SMB}}$, while equation (18) depends on $E = E_{\text{iso}}^{\text{sMB}}$ and $R_{\text{obs}}^{\text{SMB}} = R_{\text{obs}}^{\text{SMB}} f_0^{-1}$, where $f_0 < 1$ is the fraction of the 4π sr covered by the jet and $E_{\text{iso}}$ is the isotropic equivalent energy in the afterglow blast wave. $\gamma$-ray observations indicate that $\gamma$-ray emission in short GRBs is very efficient and that in general $E_{\text{iso}} \sim E_{\gamma,\text{iso}}$ (Nakar 2007). In the following discussion we assume that this is the case.

$E_{\gamma,\text{iso}}$ of short GRBs ranges at least over four orders of magnitude ($10^{49}$–$10^{53}$ erg). The rate of observed short GRBs is dominated by $10^{49}$ erg bursts, and the luminosity function can be well approximated by a power law, at least in the range $10^{49}$–$10^{51}$ erg, such that $R_{\text{obs}}^{\text{SMB}}(E) \sim 10 E_{\text{iso}}^{49} Gpc^{-3}$ yr$^{-1}$ where $\alpha \approx 0.5$–1 (Guetta & Piran 2006; Nakar et al. 2006). Plugging these into equation (18) we obtain

$$N_{\text{all-sky}}^{1.4\text{GHz}} \approx 1 f_0^{5/6} E_{\text{iso}}^{1/4} \epsilon \frac{\beta_0^{1/2}}{\epsilon_{\gamma,\text{iso}}} f_{\text{iso}}^{1/4} E_{\gamma,\text{iso}}^{3/2}.$$  (19)

This equation is similar to equation (9) of Levinson et al. (2002), with the observed luminosity function already folded in.

Narrower beamed bursts (with lower $f_0$) are more numerous and they produce less total energy per burst. The positive dependence of

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Table 3. Basic properties of radio flares from the subrelativistic dynamically ejected outflow for selected cases.

| Run | Masses | $n = 1$ cm$^{-3}$ | $n = 0.1$ cm$^{-3}$ |
|-----|--------|-----------------|-------------------|
|     |        | $F_1$ (peak)$^a$ | $t$ (peak)$^a$ | $F_1$ (peak)$^a$ | $t$ (peak)$^a$ | $F_1$ (peak)$^a$ | $t$ (peak)$^a$ | $F_1$ (peak)$^a$ | $t$ (peak)$^a$ |
|     | (M⊙)  | (mJy)           | (yr)             | (mJy)            | (yr)            | (mJy)            | (yr)            | (mJy)            | (yr)            |
| 8   | 1.4–1.2| 0.09            | 4                | 0.5              | 4               | 10              | 9               | 50              | 9               |
| 12$^b$ | 1.4–1.4| 0.04            | 1.5              | 0.2              | 2               | 5               | 3               | 30              | 3               |
| 15$^c$ | 1.4–2.0| 0.3             | 5                | 2                | 6               | 50              | 10              | 200             | 10              |
| 23$^d$ | 1.4–10 | 1.5             | 4                | 4                | 8               | 200             | 10              | 1000            | 10              |

$^a$This is the peak of the subrelativistic outflow. A mildly relativistic outflow, not calculated here, may produce a stronger and earlier peak.

$^b$The canonical ns² case.

$^c$This is the maximal signal from our runs of ns² mergers.

$^d$nsbh merger.

Table 4. Properties and detection horizons (neglecting cosmological corrections) of the subrelativistic dynamically ejected outflow from 1.4–1.4 M⊙ ns² to 1.4–10 M⊙ nsbh mergers with different radio facilities.

| Radio facility | Obs freq. (GHz) | Field of view (deg²) | 1 h rms (mJy) | ns² 1 h horizon (Mpc) | ns² 10 h horizon (Mpc) | nsbh 1 h horizon (Mpc) | nsbh 10 h horizon (Mpc) |
|----------------|----------------|----------------------|---------------|-----------------------|------------------------|------------------------|------------------------|
| EVLA$^e$       | 1.4            | 0.25                 | 7             | 360 Mpc               | 200 Mpc                | 1.8 Gpc                | 1.4 Gpc                |
| ASKAP$^f$      | 1.4            | 30                   | 30            | 170 Mpc               | 100 Mpc                | 850 Mpc                | 700 Mpc                |
| MeerKAT$^g$    | 1.4            | 1.5                  | 35            | 160 Mpc               | 90 Mpc                 | 800 Mpc                | 650 Mpc                |
| Apertif$^h$    | 1.4            | 8                    | 50            | 135 Mpc               | 75 Mpc                 | 670 Mpc                | 550 Mpc                |
| LOFAR$^i$      | 0.15           | 20                   | 1000          | 70 Mpc                | 40 Mpc                 | 300 Mpc                | 250 Mpc                |

$^e$The distance at which the observed peak flux is four times the 1 h rms.

$^f$The distance at which the observed peak flux is four times the 10 h rms.

$^g$http://www.aoc.nrao.edu/evla/

$^h$http://www.astron.nl/general/apertif/apertif

$^i$http://www.atnf.csiro.au/projects/askap/technology.html

$^j$http://www.ska.ac.za/meerkat/

$^k$http://lofar.org

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8 http://www.atnf.csiro.au/projects/askap/ssps.html
equation (19) on $f_b$ implies that overall the lower energy is ‘winning’ over the increased rate, and the detectability of narrower bursts is lower. Using, equation (19) we can put a robust upper limit on the orphans rate since all the parameters are rather well constrained by observations, with the exception of $f_b$ which is $<1$ by definition. Therefore, assuming that short GRBs are beamed, the detection of the common $\sim 10^{50}$ erg bursts in a blind survey, even with next generation radio facilities, is unlikely (Nakar 2007). However, brighter events should be detectable. If the beaming is energy independent, detectability increases with the burst energy. The luminosity function possibly breaks around $10^{51}$ erg, in which case the orphans number is dominated by $10^{51}$ erg bursts. For $f_b^{-1} = 30$ we expect from these bursts $\sim 10$ orphan afterglows brighter than 0.1 mJy at a single 1.4 GHz whole sky snapshot.

So far we discussed detectability in a blind survey. A follow-up dedicated search would be, of course, more sensitive. If compact binary mergers produce short GRBs then the energy of most GW detected bursts will be faint with $E_{gw, mm} \sim 10^{50}$ erg. The chance to detect their orphan afterglows again depended on their total energy and thus on $f_b$. Equation (11) shows that if $f_b^{-1} = 30$ then detection should be difficult but possible in a dedicated search mode. Note that since the energy of the burst is low, the radio emission will evolve quickly, reaching a peak and decaying on a week time-scale, so a prompt and rather deep search will be needed.

4.3 Identification and contamination

A key issue with the detection of compact binary merger remnants in blind surveys is their identification. Frail et al. (2012) and Ofek et al. (2011) present a census of the transient radio sky. Luckily the transient radio sky at 1.4 GHz is relatively quiet. The main contamination source are radio active galactic nuclei (AGN), however, their persistent emission is typically detectable in other wavelength and/or deeper radio observations. Moreover, the signal from a compact binary merger is expected to be located within its host galaxy (otherwise the density is too low), but away from its centre. The host and the burst location within it, should be easily detectable at the relevant distances.

The only known, and guaranteed, transient 1.4 GHz source with similar properties are radio SNe. Among these typical radio SNe are the most abundant. A transient search over 1/17 of the sky with $F_{lim} = 6$ mJy at 1.4 GHz (Levinson et al. 2002; Gal-Yam et al. 2006; Ofek et al. 2010) finds one radio SN. This rate translates to $10^{-5} - 10^{-4}$ SNe in a whole sky $F_{lim} = 0.1$ mJy survey. These contaminants can be filtered in three ways. First, by detection of the SN optical light. However, the optical signal may be missed if it is heavily extincted, and given the large number of radio SNe, misidentifying even a small fraction of them may render the survey useless for our purpose. The second filter is the optically thick spectrum at high radio frequency ($\sim 10$ GHz) at early times, which is a result of the blast wave propagation in a wind. Thus, a multiwavelength radio survey can identify radio SNe. The last filter is the luminosity-time-scale relation of typical radio SNe that is induced by the outflow velocity (e.g. fig. 2 in Chevalier, Fransson & Nymark 2006). Type II SN outflows are slow, $\sim 0.1c$, and therefore their radio emission is longer/fainter than that expected for merger remnants. The common type of Ib/c radio SNe is produced by $\sim 0.2c$ blast waves but with much less energy than what we expect from a binary merger outflow, and therefore their radio emission is much fainter. The combination of any two of these filters will hopefully be enough to identify all the typical radio SNe.

Slightly different contaminators are GRBs associated SNe. Their outflows are as fast and as energetic as those that we expect from a binary merger and therefore their radio signature is similar in time-scales and luminosities. SN 1998bw-like events are detectable by a 0.1 mJy survey at 1.4 GHz up to a distance of several hundred Mpc for 40 d and their rate is 40–700 Gpc$^{-3}$ yr$^{-1}$ (Soderberg et al. 2006), implying at least several sources at any whole sky snapshot. Here only the first filter (SN optical light) and possibly the second (optically thick spectrum) can be applied. However, given the high optical luminosity of GRBs associated SNe and their relatively low number this should be enough in order to filter them out. These contaminators highlight the importance of a multiwavelength strategy where an optical survey accompanies the radio survey to best utilize both surveys’ detections.

Finally, radio is the place to look for blast waves in tenuous mediums, regardless of their origin. Any source of such an explosion, be it a binary merger, a GRB or a SN, produces a radio signature. Therefore, all the strong explosions may be detectable in a deep radio survey, this includes for example long GRB on-axis and off-axis afterglows and giant flares from extragalactic soft gamma-repeaters. The difference between the radio signatures of the different sources (amplitude, spectrum and time evolution) depends on the blast wave energy and velocity and on the external medium properties. We thus will be able to identify the characteristics of binary mergers outflows. If, however, there is a different source of $\sim 10^{50}$ erg of sub- or mildly relativistic outflow that explodes in the ISM it will be indistinguishable from binary mergers (at least in the radio). Currently we are not aware of any such source, with the exception of long GRBs at the low end of the luminosity function, but these are too rare to contaminate a survey. Any other source of such outflows, if existent, will probably be a part of the family of collapsing/coalescing compact objects.

5 IR–UV TRANSIENTS FROM RADIOACTIVE DECAYS, ‘MACRONOVA’

The ejected material is extremely neutron rich and rapidly expanding. Under such conditions rapid neutron capture is hard to avoid (Hoffman et al. 1996; Freiburghaus, Rosswog & Thielemann 1999; Roberts et al. 2011; Korobkin et al. 2012). The latter study finds that ns$^2$ and mshb mergers produce a unique, solar-system-like $r$-process abundance pattern for nucleon numbers $A > 120$, independent of the astrophysical parameters of the merging binary system. The $r$-process itself occurs on a very short time-scale, but the freshly synthesized elements subsequently undergo nuclear fission, $\alpha$- and $\beta$-decay which occur on much longer time-scales. The SN-like emission from this expanding material was first suggested by Li & Paczyński (1998) and discussed later by Kulkarni (2005) and Metzger et al. (2010). We combine the ejecta velocity profiles found in our simulations with the time-dependent radioactive power injection found in the recent $r$-process study of Korobkin et al. (2012) to calculate bolometric light curves.

The optical depth of the expanding outflow decreases with time, and as a result larger amounts of mass become visible to the observer. The ejected mass has a gradient of velocities. We denote as $m(v)$ the mass with asymptotic velocity at infinity $\sim v$. Before the peak of the emission, the observed mass, $m_{obs}$, is the one for which the diffusion time is comparable to the expansion time, namely $\tau_{obs} \approx c/v$, where $c(m) \approx \kappa m/(4\pi v^2)$ (note that $m = m(v)$) and $\kappa$ is the opacity cross-section per unit of mass. The emission peaks at $t_{peak}$, when the entire ejecta become observable, namely at the first
time that $m_{\text{obs}} = m_{\text{ej}}$. At later time the entire ejecta are exposed and $m_{\text{obs}}(t > t_{\text{peak}}) = m_{\text{ej}}$.

A major uncertainty in the light-curve calculation is the opacity of r-process elements. The opacity determines first the bolometric luminosity and second its spectrum. Here we assume a constant and grey opacity in order to determine the luminosity, which is less sensitive to the detailed wavelength dependence of the opacity. We do not attempt to predict the observed spectrum. Metzger et al. (2010) discussed the opacity of the neutron-rich ejecta expect it to be similar to that of iron-group elements, which they approximate as a constant $\kappa = 0.1 \text{ cm}^2 \text{ g}^{-1}$. Below use it as the canonical value.

At any time prior to $t_{\text{peak}}$ the observed mass is

$$m_{\text{obs}}(t < t_{\text{peak}}) \approx 0.05 M_\odot \left( \frac{\kappa}{0.1 \text{ cm}^2 \text{ g}^{-1}} \right)^{-1} \left( \frac{v}{0.1 c} \right)^2 \left( \frac{t}{t_{\text{d}}} \right)^2.$$  \hspace{1cm} (20)

Note that this is an implicit equation since $v$ itself depends on $m_{\text{obs}}$. Thus, $m_{\text{obs}}$ increases with time (slightly slower than $t^2$, since $v$ decreases with $m$) until all of the ejecta is exposed at $t_{\text{peak}}$. The observed luminosity is dominated by the energy release via radioactive decay. Korobkin et al. (2012) follow the nucleosynthesis of a number of fluid elements that are ejected at different velocities and during different stages of the merger and find that all of them result in a similar radioactive energy injection rate per unit of mass into the expanding ejecta that can be approximated by

$$\dot{e} = 2 \times 10^{18} \left[ \frac{1}{2} - \frac{1}{\pi} \arctan \left( \frac{t - t_0}{\sigma} \right) \right]^{1.3} \frac{\epsilon_{\text{therm}}}{0.5} \text{ erg g}^{-1} \text{ s}^{-1},$$  \hspace{1cm} (21)

where $t_0 = 1.3$ s and $\sigma = 0.11$ s and $\epsilon_{\text{therm}}$ is the fraction of injected energy that is emitted as thermalized radiation and not as unobserved neutrinos or $\gamma$-rays. The value of $\epsilon_{\text{therm}}$ is not well constrained and below we use the default value of Korobkin et al. (2012), $\epsilon_{\text{therm}} = 0.5$. On the relevant time-scales equation (21) corresponds to the Li & Paczyński (1998) coefficient $f = 10^{-6} (t/d)^{-0.3}$. The resulting bolometric (thermalized) luminosity is

$$L \sim \dot{m} \dot{e} \sim 2 \times 10^{41} \text{ erg s}^{-1} \left( \frac{\epsilon_{\text{therm}}}{0.5} \right) \left( \frac{m_{\text{obs}}}{10^{-2} M_\odot} \right) \left( \frac{t}{t_{\text{d}}} \right)^{-1.3}.$$  \hspace{1cm} (22)

Note that at early times $L$ increases because $m_{\text{obs}}$ increases almost like $t^2$. This luminosity is emitted in the IR to UV range. At early time $L$ is increasing (slightly slower than linearly). The luminosity peaks when the entire ejecta is seen (roughly at $t = 0.5$ d), and the light-curve decays afterwards roughly as $t^{-1.3}$. The value of $m_{\text{obs}}$ is inversely proportional to the opacity. Therefore higher opacity implies a slower evolving and dimmer signal, and it most likely results in a redder transient.

The resulting macronova light curves, assuming $\kappa = 0.1 \text{ cm}^2 \text{ g}^{-1}$, are shown in Fig. 13 for different ns$^2$ mergers and the basic properties are summarized in Table 5. The canonical ns$^2$ merger case (thick black line) peaks after 0.4 d with $5 \times 10^{41}$ erg s$^{-1}$. The macronova light curves produced by the two nsbh mergers that we consider are depicted in Fig. 14, the canonical ns$^2$ case ($2 \times 1.4 M_\odot$) is shown for reference purposes. Since more mass is ejected these macronova peak later, at about 0.7 d and their peak luminosities can reach about $10^{42}$ erg s$^{-1}$. Our estimates of the macronova peak luminosity for the canonical ns$^2$ is similar to the calculation of Metzger et al. (2010) for $m_{\text{ej}} = 0.01 M_\odot$ and $v = 0.1c$. Other merger cases are brighter by a factor of 2–3 mostly due to the larger amount of ejected mass. Fig. 15 shows the effect of the

![Figure 13. Bolometric light curves of 'macronovae' – SN-like event that are powered by a radioactive decays within the dynamical ejecta. The opacity is assumed to be grey and constant with $\kappa = 0.1 \text{ cm}^2 \text{ g}^{-1}$. Most of the luminosity is emitted in IR, optical and UV. The canonical ns$^2$ case.](https://example.com/figure13)

![Figure 14. Bolometric light curves of 'macronovae' – SN-like event that are powered by a radioactive decays within the dynamical ejecta. The opacity is assumed to be grey and constant with $\kappa = 0.1 \text{ cm}^2 \text{ g}^{-1}$. Most of the luminosity is emitted in IR, optical and UV. The light curves are for nsbh mergers of different masses. For a comparison the light curve for the canonical 1.4–1.4 $M_\odot$ merger is also shown.](https://example.com/figure14)
6 CONCLUSION

Compact binary mergers are expected to eject subrelativistic, mildly relativistic and possibly ultrarelativistic outflows. The interaction of such an outflow with the circummerger environment produces a long-lived radio remnant. Radioactive decay within the outflow produces a short-lived IR–UV transient.

To estimate the properties of these transients we have carried out a series of merger simulations which aims at finding the properties of the dynamically ejected mass. This is a lower limit on the amount of the ejected mass, since other sources of sub-, mildly and ultrarelativistic outflows are not accounted for in our simulations, see the discussion at the end of Section 2. We find that our canonical merger case of two 1.4 M⊙ ns dynamically ejects 0.013 M⊙ with a distribution of velocities in the range of 0.05–0.2c and an average value \( v \approx 0.1c \). The energy carried by this outflow is \( 1.6 \times 10^{50} \text{erg} \). Other ns\(^2\) mergers, and especially those with unequal masses, generate more massive outflows at slightly faster velocities, up to 0.04 M⊙ and 9 \( \times 10^{50} \text{erg} \) of kinetic energy in some cases. The dynamically ejected outflow that we find for nsbh mergers carries about 10\(^{51} \text{erg} \).

A strong radio remnant is expected in any merger of a ns\(^2\) whose properties are similar to those of the Galactic ns binary population (i.e. typical circummerger Galactic disc density of 1 cm\(^{-3}\)). Such a radio remnant appears months to years after the merger and remains bright for a similar time. Therefore, a trigger following the detection of a GW signal can wait weeks after the event and no online triggering is needed. In addition the long lifetime of the remnants enables their detection in a blind survey. Here we calculate the light curves from 1 yr after the merger, since our simulations do not include the mildly relativistic outflow which dominates the emission at earlier times. The radio flux from a canonical 1.4–1.4 M⊙ ns\(^2\) merger taking place at the advanced LIGO/Virgo horizon for such mergers, 300 Mpc, in \( n = 1 \text{ cm}^{-3} \) environment is 0.04 mJy (0.2 mJy) at 1.4 GHz (150 MHz) 1–4 yr after the merger. This signal could be easily detected by a 1 h observation of the EVLA or by a whole day observation on ASKAP or MeerKAT. The sensitivity of present lower frequency e.g., at 150 MHz, is insufficient for a detection at the advanced LIGO/Virgo horizon. Longer observations (e.g. 10 h on the EVLA) can detect these mergers even if they are in a lower density environment (\( n = 0.1 \text{ cm}^{-3} \)). A mildly relativistic component in the ejecta probably increases the brightness and detectability of the signal on time-scales of weeks to months. The nsbh GW horizon is farther than the ns\(^2\) GW horizon. Our numerical simulations find that nsbh mergers produce higher energy outflows resulting in larger fluxes. Overall we find that the ELVA detection horizon for nsbh is almost a factor of 2 larger than the advanced LIGO/Virgo detection horizon.

We find that the optimal frequency to carry out a search for merger remnants is 1.4 GHz. Taking a subrelativistic outflow with 10\(^{50} \text{erg} \) and a canonical ns\(^2\) merger rate of 300 Gpc\(^{-3}\) yr\(^{-1}\)(and a range of 20–2 \( \times 10^3 \text{ Gpc}^{-3}\) yr\(^{-1}\)) we expect a detection of \( \sim 20 \) (1–1200 correspondingly) radio ns\(^2\) remnants in a 0.1 mJy all sky survey. The expected higher velocity component increases this rate making remnants detectable even in a survey that covers only a small fraction of the sky or that operates at a mJy sensitivity. Thus, a sensitive large field-of-view GHz survey by currently available facilities has a great potential to constrain the rate of binary mergers, information that is of great importance for the design and operation of the advanced GW detectors.

The detectability of the radio transients depends strongly on the circummerger density, which may be low if the binary has been ejected from its host galaxy before the merger. This uncertainty implies that there may be a fraction of ns\(^2\) mergers, that take place out of the disc of Milky-Way-like galaxies, whose radio flares are faint. However our estimates of the detection rate in a blind radio survey are less affected by this uncertainty as our canonical circummerger density and the expected merger rate are based on the observed Galactic ns\(^2\) binaries, which are all located in the Galactic disc. If there is a population of mergers which take place out of the disc of Milky-Way-like galaxies, it will be in addition to the population that we consider here for radio transient blind searches.

These radio transients should be compared with short GRB orphan radio afterglows. These may be produced by compact binary
mergers if they are launching also ultrarelativistic outflows. Our estimates of orphan afterglows detectability are based on short GRB observations and are therefore independent of whether short GRBs are binary mergers or not. The main uncertainty in the rate estimates is the GRB beaming factor. We find that assuming $f_{\text{beam}} = 30$ there are about a dozen orphan afterglows at 0.1 mJy in a single 1.4 GHz whole sky snapshot. These are dominated by relatively energetic short GRBs ($E_{\gamma,\text{iso}} \sim 10^{51}$) and their duration is several weeks. If binary mergers are short GRB engines then a GW-triggered event will most likely be of a low-energy GRBs, $E_{\gamma,\text{iso}} \sim 10^{50}$ erg, and a true energy, after beaming correction, that is even lower. Their radio orphan afterglow will probably still be detectable in a deep search. However, its variability time-scale is short, about a week, so the search should be done promptly.

Our results show the great potential of 1.4 GHz radio transient observations at the sub-mJy level for the detection of ns mergers. On the observational side these predictions provide an excellent motivation for carrying out a whole sky sub-mJy survey using the EVLA or other upcoming radio telescopes. The main source of contamination in such surveys would be radio SNe and those could be distinguished from compact binary mergers by their optical signal, optically thick spectrum and other characteristic properties.

The IR–UV light curves expected from ‘macronovae’, SN-like events powered by the radioactive decay within the ejecta, depends on the total mass ejected and on the velocity distribution. We present a general method for performing such calculations. For each of the 23 simulations that cover the binary parameter space we follow a large number of ejecta trajectories with a state-of-the-art nuclear reaction network. All trajectories show very similar radioactive heating histories which were fitted in Korobkin et al. (2012), see their equation (12). A large uncertainty, though, comes from the poorly known outflow opacity. If we adopt the value that has been discussed in some detail in Metzger et al. (2010), we find that ns mergers peak at about $5 \times 10^{41} \text{erg s}^{-1}$, corresponding to absolute bolometric magnitudes of $-15 \rightarrow -16$. The magnitude in a given observed optical band is probably fainter by a magnitude or two. Such events can be detected by current blind surveys like PTF up to a distance of about 150 [300] Mpc and by LSST up to 0.8 [1.5] Gpc. The short duration of these events, about 0.4 to 0.7 d, may pose problem as it would require very short cadence surveys. Factoring in these limits we expect for the canonical rate of 300 Gpc$^{-3}$ yr$^{-1}$ ns mergers a detection of 0.01 macronovae yr$^{-1}$ by PTF and 100 yr$^{-1}$ by the LSST. If the opacity is higher than the value suggested by Metzger et al. (2010) then the peak time is delayed while the peak luminosity drops, making the detection of macronova light even harder.

Before concluding we address the relation between mergers the associated radio flares and the short GRBs. While it has not been confirmed it is possible, and maybe even likely, that compact binary mergers are the origin of short GRBs. For that reason we compared ‘orphan afterglows’ with radio flares and demonstrated that radio flares are expected to be brighter. It has been claimed (Metzger & Berger 2012) that short GRBs arise in large distances from the host galaxies regions and that low density is inferred from modelling of their afterglow. Hence they suggest that if the association of mergers and GRBs is correct this implies that mergers take place in low-density regions and as such their radio flares will be undetectable. However, we know that regardless of the question whether short GRBs are associated with mergers or not, the compact binaries that have been observed in our galaxy are in the galactic disc, namely in high ISM density. Merger rate estimates based on these binaries (Kaloger et al. 2004a,b) suggest a comparable rate of events to the beaming corrected rate inferred from short GRBs. This implies that, while there may (or may not) be a merger population at low-density environment, there must be (regardless of the whether there is a connection to short GRBs) a large population of the mergers that take place in Milky Way ISM-like density.

We have presented here the detailed methodology for calculation of the two more robust EM transients that should accompany compact binary mergers – radio flares that arise from the interaction of sub- mildly relativistic outflows with the surrounding matter and macronovae that arise from the radioactive decay of the ns matter. We have obtained these estimates from realistic (employing a slew of microphysics ingredients) though Newtonian merger simulations. The major sources of uncertainty in our estimates for radio flares are the fraction of mildly relativistic ejecta (which can only increase the emission) and the surrounding matter density. For the IR–UV macronovae the main uncertainties are the radioactive energy source within the ns matter and its opacity. We find that the prospects for detection of both radio flares and the IR–UV macronova are promising by intensive follow up searches following GW signals. Radio flares that last a few months to years have the advantage that they do not require a rapid follow-up and that the background sky contamination is rather low. Macronovae are more challenging, as they require a large field of view follow-up within a very short time frame of days or even less and they need to be identified in the crowded optical transient sky. Their advantage is that their emission is independent of the circummerger environment.

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