Detectability of compact binary merger macronovae

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Abstract

We study the optical and near-infrared luminosities and detectability of radioactively powered electromagnetic transients (‘macronovae’) occurring in the aftermath of binary neutron star and neutron star black hole mergers. We explore the transients that result from the dynamic ejecta and those from different types of wind outflows. Based on full nuclear network simulations we calculate the resulting light curves in different wavelength bands. We scrutinize the robustness of the results by comparing (a) two different nuclear reaction networks and (b) two macronova models. We explore in particular how sensitive the results are to the production of α-decaying trans-lead nuclei. We compare two frequently used mass models: the finite-range Droplet model (FRDM) and the nuclear mass model of Duflo and Zuker (DZ31).

We find that the abundance of α-decaying trans-lead nuclei has a significant impact on the observability of the resulting macronovae. For example, the DZ31 model yields considerably larger abundances resulting in larger heating rates and thermalization efficiencies and therefore predicts substantially brighter macronova transients. We find that the dynamic ejecta from NSNS models can reach peak K-band magnitudes in excess of $-15$ while those from NSBH cases can reach beyond $-16$. Similar values can be reached by
some of our wind models. Several of our models (both wind and dynamic ejecta) yield properties that are similar to the transient that was observed in the aftermath of the short GRB 130603B. We further explore the expected macronova detection frequencies for current and future instruments such as VISTA, ZTF and LSST.

Keywords: electromagnetic transients, gravitational waves, neutron stars, nucleosynthesis, accretion

(Some figures may appear in colour only in the online journal)

1. Introduction

The recent, direct detection of gravitational waves [1] has finally opened the door to the long-awaited era of gravitational wave (GW) astronomy. With a sky localization uncertainty of $\sim 600$ square degrees for the first event, however, the astronomical environment (e.g. type of galaxy or ambient medium density) in which the black hole (BH) merger took place is essentially unknown. To extract this information, measure the event’s redshift and to constrain the evolutionary channels that lead to the merger in the first place one needs a coincident electromagnetic (EM) signal.

While today the majority of LIGO-detectable sources is believed to be binary black holes, advanced detectors should also be able to observe compact binary systems [2] that contain at least one neutron star (NS; either NSNS or NSBH; hereafter collectively referred to as compact binary mergers). In such mergers neutron star matter is decompressed and ejected into space. Within the ejecta rapid neutron capture ($r$-process) synthesizes a range of heavy elements [3–7] up to and beyond the platinum peak near $A = 195$. The subsequent radioactive decay of the freshly synthesized $r$-process elements in the expanding ejecta causes EM transients known as ‘macronovae’ (MN) [9–14]. They are promising EM counterparts of GWs [15], since they are—contrary to gamma-ray bursts—close to isotropic and viewing angle effects are only of order unity [16, 17]. By now, there are several cases where near-infrared excesses in the aftermath of gamma-ray bursts have been interpreted as being due to macronovae [18–22].

Compact binary mergers eject matter via several channels, see figure 1 and e.g. [28] and [29] for recent reviews. The best-studied channel (stage ‘II’ in figure 1) is the so-called dynamic ejecta [6, 23–25, 30–33] that become unbound immediately at merger. They are launched by either gravitational torques (‘tidal component’) or, in the NSNS case, by hydrodynamic interaction at the interface between the stars (‘interaction component’, see [31] and [34]). While the tidal component carries away matter with the original, very low electron fraction ($Y_e \approx 0.04$) which is set by cold $\beta$-equilibrium in the original neutron star, the interaction component is heated and can increase its $Y_e$ via positron captures $e^+ + n \rightarrow p + \bar{\nu}_e$.

Both sub-components, however, share the property that they are hardly impacted by neutrino irradiation, see e.g. [35], simply because they have already reached large distances from the remnant before the neutrino luminosity starts rising in earnest (after accretion torus formation at $t > 10\,\text{ms}$ after contact). Nevertheless, for material that is not immediately ejected, weak

$^7$The same phenomenon is often referred to as ‘kilonova’. We prefer ‘macronova’ (MN), since the phenomenon is not a thousand times brighter than a nova, as originally thought. See [8] for discussion of the naming conventions.

$^8$See figure 2 in the latter publication for an illustration.
interactions can change the electron fraction substantially, see e.g. [36] who find a particularly large fraction of high $Y_e$ material in the ejecta.

The dynamically ejected matter is complemented by baryonic winds that are driven by neutrino-energy deposition, magnetic fields, viscous evolution and/or nuclear recombination energy [37–46]. If a completely or temporarily stable neutron star survives in the centre of a merger remnant (see stage ‘III’ in figure 1) the neutrino- and the magnetically driven winds are substantially enhanced in comparison to the case where a BH forms immediately. Simulations indicate that a prompt collapse of the massive neutron star to a black hole can be avoided in many cases [47–51] even if the central mass substantially exceeds the Tolman-Oppenheimer-Volkoff (TOV) limit. The calculations of [52] and [53, 54] indicate that prompt collapse is avoided, unless the initial mass of the binary exceeds the TOV mass limit by more than $\sim 35\%$. The two precisely determined neutron star masses near 2.0 $\odot$ (J1614-2230, [55] and PSR J0348+0342, [56]), place this threshold value to at least 2.7 $\odot$. Thus, a large fraction of the mergers may go through such a metastable phase with a strong neutrino-driven wind phase and binaries at the low mass end may even produce stable, very massive neutron stars.

Recent studies [57, 58] indicate, however, that a BH needs to form on a time scale of less than $\sim 100\,\text{ms}$ in order to launch a short GRB (stage ‘V’ in figure 1), otherwise outflows will be baryon-overloaded and will not reach relativistic speeds. Seconds after the burst the resulting BH-disk system can release a good fraction of its disk (stage ‘VI’). Recent work [40, 43] suggests that neutrinos play a sub-dominant role for the mass loss once a BH has formed.

NSBH mergers (stage ‘VII’) can release substantially more matter in dynamic ejecta than NSNS binaries [11, 26, 27, 59–62]. If a sufficiently massive torus forms around the BH, it is expected that a jet can be launched into a relatively unpolluted surrounding (stage ‘VIII’). But large BH-masses in combination with low spins may lead to very small or no disks [29, 63]. On longer times scales (seconds) a fair fraction of the initial disk mass is expected to become unbound (stage ‘IX’), just as in the NSNS case.
Recently, a fast neutron component has been discussed to produce a macronova precursor signal [64]. This model is based on SPH simulations in the conformal flatness approximation where some particles are ejected early on with a high velocity from the shear interface between the two merging neutron stars. While such a precursor is a possibility, we do not see it in our simulations and neither do, e.g. [32] in their recent GR simulations. Therefore, we will not consider this possibility in the following discussion.

Our knowledge of the rates of compact binary mergers is still plagued by large uncertainties. Simple constraints can be derived by assuming that they are related to the production of r-process elements. If we take the solar-system r-process abundance pattern [65] as representative and define a quantity \( \sigma = \sum_{A>i} X_A^{\text{r}} \), where \( X_A^{\text{r}} \) is the r-process mass fraction with nucleon number \( A \), and use a Milky Way baryonic mass of \( M_{\text{MW}} = 6 \times 10^{11} M_\odot \) [66], we find that the Milky Way contains \( \approx 19000 \, M_\odot \) of r-process material in total. About \( M_{A>130} = \sigma(130)M_{\text{MW}} = 2530 \, M_\odot \) of this matter has \( A > 130 \) and about 500 \( M_\odot \) are beyond the ‘platinum peak’ (\( A \approx 195 \)). With an age of the Galaxy of \( t_{\text{MW}} \approx 10^{10} \, \text{yrs} \), this yields average production rates of \( \dot{M}_{\text{r,all}} = 1.9 \times 10^{-6} \, M_\odot \, \text{yr}^{-1} \), \( \dot{M}_{A>130} = 2.5 \times 10^{-7} \, M_\odot \, \text{yr}^{-1} \) and \( \dot{M}_{A>195} = 5 \times 10^{-9} \, M_\odot \, \text{yr}^{-1} \). The product of average ejecta mass and event rate is known but the individual factors are not. This is shown as lines from the upper left to the lower right in figure 2 (e.g. r-process with \( A > 130 \) in red). As an example, if an event that produces all r-process \( A > 130 \) occurs at a rate of \( 10^{-5} \, \text{yr}^{-1} \) it has to eject on average \( \times 10^{-2} \, M_\odot \) per event. For comparison, some representative simulation results for both NSNS [24, 25, 67] and NSBH [26, 27] are also indicated. We also translated the rates from yr\(^{-1}\) MWEG\(^{-1}\) (bottom axis) to yr\(^{-1}\) Gpc\(^{-3}\) (axis on top) via a density of \( \times 10^{-2} \, \text{MWEG Mpc}^{-3} \) [68]. Here MWEG abbreviates ‘Milky Way equivalent galaxy’. Also indicated are the expected LIGO upper limits of science runs O1-O3 [69]. We have further marked the NSNS merger rates from the population synthesis studies of [70] and the range of sGRB rates estimated by [71].

The paper is organized as follows. Our study is based on a multitude of ingredients including hydrodynamic and nuclear reaction network simulations, macronova emission models and tools to explore the detectability of macronovae with existing and future instruments. The methodological elements that enter our study are described in section 2.2. For our simulations of NSNS mergers we use a multi-physics numerical model that couples Newtonian, ideal hydrodynamics discretized via the Smoothed Particle Hydrodynamics method with a nuclear equation of state and neutrino reactions. The neutrino-emitting weak interactions cool the merger remnant and the corresponding weak interactions can change the neutron to proton ratio, a quantity that is crucial for the subsequent nucleosynthesis and the resulting electromagnetic emission. We further discuss in this section the nuclear reaction networks that we use and our macronova models that receive their energy input directly from the nuclear networks. Our baseline model is the one of [17], which we extend to a second model that includes time-dependent thermalization efficiencies and two different nuclear mass formulae. Section 3 presents our results on nucleosynthesis and discusses the impact of different nuclear mass formulae on the resulting abundances, the thermalization efficiency and the radioactive heating rates. We further present optical and near-infrared lightcurves and we discuss the detection feasibility, the follow-up of LIGO triggers and the use of GRB-triggers to search for macronovae. Section 4 summarizes the main findings of this study.

2. Methodology

In this study we explore macronova transients based on nuclear network calculations along thermodynamic trajectories. To obtain these trajectories, we perform a set of hydrodynamic
simulations for dynamic NSNS ejecta, for other cases we use a parametrized treatment with numerical values based on existing hydrodynamic studies.

2.1. NSNS merger simulations

The NSNS simulations of this paper make use of the Smooth Particle Hydrodynamics (SPH) method, see [72–75] for recent reviews. Our code is an updated version of the one that was used in earlier studies [11, 76–78]. We solve the Newtonian, ideal hydrodynamics equations for each particle $a$:

\[
\frac{d\mathbf{v}_a}{dt} = - \sum_b m_b \left( \frac{P_{ab}}{\rho_a} + \frac{P_{ba}}{\rho_b} + \Pi_{ab} \right) \nabla_{ab} W_{ab} + f_{a,G} + f_{a,GW} \\
\frac{du_a}{dt} = \sum_b m_b \left( \frac{P_{ab}}{\rho_a} + \frac{1}{2} \Pi_{ab} \right) \mathbf{v}_{ab} \cdot \nabla_{ab} W_{ab} - \frac{du_a'}{dt}
\]  

Figure 2. Summary of various rate constraints. The lines from the upper left to lower right indicate the typical ejecta mass required to explain all r-process/all r-process with $A > 80$/all r-process with $A > 130$ for a given event rate (lower panel per year and Milky Way-type galaxy, upper panel per year and Gpc$^3$). Also marked is the compiled rate range from Abadie et al (2010) for both double neutron stars and neutron star black hole systems and (expected) LIGO upper limits for O1 to O3 (Abbott et al 2016b). The dynamic ejecta results from some hydrodynamic simulations are also indicated: the double arrow denoted ‘nsns Bauswein + 13’ indicates the ejecta mass range found in [23], ‘nsns Rosswog 13’ refers to [24], ‘nsns Hotokezaka + 13’ to [25], ‘nsbh Foucart + 14’ to [26] and ‘nsbh Kyutoku + 13’ to [27].
the mass density is calculated by summing up contributions from neighboring particles
\[ \rho_b = \sum_b m_b W_{ab}. \] (4)

Here \( m_b \) is the mass of particle \( b \) and \( W_{ab} = W(\bar{r}_{ab} - \bar{r}_b, h_{ab}) \) denotes a suitably chosen smoothing kernel [79] evaluated with the average smoothing length \( h_{ab} = (h_a + h_b)/2 \). For the kernel function, we choose the Wendland C6 kernel [80, 81] together with 200 neighbours per particle. We have recently scrutinized various ingredients in the SPH method [82] and found that this kernel has excellent numerical properties and in particular drastically reduces the noise in an SPH simulation in comparison to the cubic spline kernel that is commonly used. \( \bar{f}_{ab; GW} \) is the additional acceleration due to the Newtonian self-gravity of the fluid that we calculate using the binary tree of [83] and \( \bar{f}_{h; GW} \) is a simple prescription for gravitational wave backreaction force [84].

The relative particle velocity is denoted as \( \bar{v}_{ab} = \bar{v}_a - \bar{v}_b \). To produce entropy in shocks, artificial dissipation is included via the tensor \( \Pi_{ab} \). It has the standard form described in [85] but particular care has been taken to avoid possible artifacts due to artificial viscosity, this has been outlined in detail in [86]. The quantity \( u_a \) denotes the specific internal energy of particle \( a \) the evolution of which is determined by \( PdV \)-work and viscous heating (summation term) and by the energy loss to neutrinos, \( \frac{\partial e}{\partial t} \). The quantities
\[ \lambda_{PC} = \frac{R_{PC}^{\text{eff}}}{\eta_{np}} \quad \text{and} \quad \lambda_{EC} = \frac{R_{EC}^{\text{eff}}}{\eta_{pn}} \] (5)

are the electron and positron capture rates per neutron/proton. \( R_{EC/PC}^{\text{eff}} \) are the effective neutrino number emission rates and the quantities \( \eta_{np}/\eta_{pn} \) reduce in the non-degenerate limit to the number densities \( n_n \) and \( n_p \) [87], for a more detailed account on our opacity-dependent multi-flavour leakage scheme, we refer to our original paper [77]. Our treatment includes in particular electron and positron captures and therefore the nuclear matter can change its electron fraction \( Y_e \) in the course of the merger. The pressure at a particle \( b \), \( P_b(\rho_b, T_b, Y_e, b) \), is calculated using the temperature-dependent relativistic mean field equation of state (EOS) of Shen et al [88, 89], extended to lower densities as described in [76]. The binary systems that we explore in this study are summarized in table 1, each of them is modeled with \( 10^6 \) SPH particles and the stars have negligible initial spin as expected due to the low viscosity of neutron star matter and the very short tidal interaction time [90, 91]. The initial electron fraction within the neutron stars is given by the cold, \( \beta \)-equilibrium condition calculated from the chemical potentials of our nuclear EOS. Our typical simulation time is \( \sim 30 \) ms. As an example, a series of snapshots from a NSNS simulation (1.3–1.3 \( M_\odot \); run N2; color-coded is the electron fraction \( Y_e \)) is shown in figure 3.

2.2. Ejecta

We identify unbound matter in our hydrodynamic simulations by the criterion \( \sqrt{\frac{v^2}{2} + \phi} > 0 \) at the end of our simulation, where \( \alpha \) labels the SPH particle and \( v \) and \( \phi \) are velocity and gravitational potential. To double-check this criterion we compare with a criterion based on the outward radial velocity being larger than the local escape velocity. The ejecta masses based on both criteria agree with each other to within \( \sim 2\% \). To have a robust upper limit, we also
The result of the simulations, $t_{\text{end}}$, in ms and the average velocities at infinity, $(v_{\infty,j})$, in units of $c$. The quantity $\chi$ for the NSBH cases refers to the dimensionless black hole spin parameter.

### Table 1. Overview over the simulation parameters. Also shown are the resulting mass fractions of lanthanide ($X_{\text{lan}}$) and actinide ($X_{\text{act}}$) elements. All NSBH cases start with an initial entropy of $2\,k_B/nuc$ and $Y_e = 0.06$. The masses $m_1$, $m_2$, $m_{\text{ej}}$, and $m_{\text{ej,max}}$ are given in solar units, $m_{\odot}$ and $m_{\text{ej,max}}$ in $10^{-2}\,M_{\odot}$. The end times of the simulations, $t_{\text{end}}$, in ms and the average velocities at infinity, $(v_{\infty,j})$, in units of $c$. The quantity $\chi$ for the NSBH cases refers to the dimensionless black hole spin parameter.

#### Dynamic ejecta NSNS mergers

| Run | $m_1$ | $m_2$ | $t_{\text{end}}$ | $m_{\text{ej}}$ | $(v_{\infty,j})$ | $m_{\text{ej,max}}$ | $X_{\text{lan}}$ | $X_{\text{act}}$ |
|-----|------|------|-----------------|-----------------|----------------|-----------------|----------------|----------------|
| N1  | 1.2  | 1.2  | 32.1            | 0.79            | 0.12           | 3.17            | 0.172          | 0.051          |
| N2  | 1.3  | 1.3  | 31.1            | 1.26            | 0.11           | 2.70            | 0.170          | 0.057          |
| N3  | 1.4  | 1.4  | 38.3            | 0.84            | 0.11           | 2.25            | 0.179          | 0.060          |
| N4  | 1.2  | 1.4  | 30.6            | 1.59            | 0.11           | 2.92            | 0.178          | 0.054          |
| N5  | 1.4  | 1.8  | 25.3            | 3.40            | 0.12           | 4.76            | 0.169          | 0.076          |

#### Dynamic ejecta NSBH mergers

| Run | $m_{\text{ns}}$ | $m_{\text{bh}}$ | $\chi$ | $m_{\text{ej}}$ | $(v_{\infty,j})$ | $X_{\text{lan}}$ | $X_{\text{act}}$ | Comment |
|-----|----------------|----------------|--------|----------------|----------------|----------------|----------------|---------|
| B1  | 1.4            | 7.0            | 0.7    | 4.0            | 0.20           | 0.199          | 0.041          | Foucart + (2014), run M14-7-S7 |
| B2  | 1.4            | 7.0            | 0.9    | 7.0            | 0.18           | 0.193          | 0.048          | Foucart + (2014), run M14-7-S9 |
| B3  | 1.2            | 7.0            | 0.9    | 16.0           | 0.25           | 0.195          | 0.046          | Foucart + (2014), run M14-7-S9 |

Figure 3. Electron fraction in a 1.3–1.3 $M_{\odot}$ merger (model N2; only matter below orbital plane shown) at $t = 7.06$, 11.6 and 12.4 ms.

Examine $v_j^2 + \phi_j + u_a > 0$, where $u_a$ is the specific internal energy of a particle $a$. The resulting mass is noted as $m_{\text{ej,max}}$ in table 1, and it is used to set the upper limit that can be plausibly expected for the electromagnetic signal from dynamic ejecta.

Studies that focus on the long-term evolution of accretion disks around BHs [40, 43, 92–94] find that a substantial fraction ($\sim 20\%$) of the initial torus can become unbound. According to [95] the torus masses can, depending on the initial mass ratio, be very large and reach multiples of $0.1\ M_{\odot}$. Thus the unbound mass from accretion tori can be large and actually rival the dynamic ejecta masses of even asymmetric mass ratio NSNS/NSBH mergers. The exploratory numerical studies, however, suggest that the velocities are substantially lower than in the dynamic ejecta case. [43], for example, find that the average velocities of the torus component never exceed $0.06c$, while the dynamic ejecta in the NSBH cases can be larger than $0.2c$, see table 1.

We subsume all the ejecta types other than dynamic ejecta under the broad category ‘winds’. Despite their different origin, they all have in common that they are exposed for
longer time to the neutrino irradiation, therefore this material has a larger $Y_e$ and consequently a different nucleosynthesis. In particular, such matter has a lower content in those elements that are the major opacity sources [96–99]. With potentially different nuclear heating rates and lower opacities, these winds are promising sources for EM transients, that could potentially outshine the signals from the dynamic ejecta, if they are not obscured by them. A first study by [100] finds, however, that due to their different velocities, the wind expansion at large radii is not influenced by the dynamic ejecta.

Our parametric wind model contains four parameters: the wind mass $m_w$, the initial entropy $s_0$, the electron fraction $Y_e$ and the (terminal) wind velocity $v_w$. To create initial conditions for the network and subsequent macronova calculations, we also need a starting temperature $T_0$ and an initial radius $R_0$. The starting temperature is chosen as $T_0 = 9 \times 10^9$ K, so that we are safely above the threshold temperature for nuclear statistical equilibrium (NSE, $T_{\text{NSE}} \approx 5 \times 10^9$ K). Thus, the initial abundance distribution is set by NSE and, below $8 \times 10^9$ K, the nuclear reaction network takes over. The initial radius $R_0$ is found by the requirement that the average density $\rho = \frac{3m_w}{4\pi R_0^3}$ together with $T_0$ reproduces the desired initial entropy $s_0$. To keep the parameter space under control and motivated by the strongly peaked entropy distribution of [41, 43] and [100], we fix the initial entropy to values of $15 \, \text{keV}$ per baryon. Typical values for $R_0$ are $\sim 650$ km, in reasonable agreement with the numerical studies of neutrino-driven winds [39, 41]. Since we are interested here in exploring the lower opacity case, we restrict our study to electron fractions above the threshold value for heavy r-process, $Y_{e,\text{thresh}} = 0.25$, see figure 8 in [34]. From the results shown in [41], we expect terminal wind velocities around $0.05c$. These numbers motivate our ‘wind ’ simulation. Once the parameters have been set, the density evolves according to $\rho_0 = \frac{3m_w}{(4\pi R_0^3)}$ together with $T_0$ reproduces the desired initial entropy $s_0$. To keep the parameter space under control and motivated by the strongly peaked entropy distribution of [41, 43] and [100], we fix the initial entropy to values of $15 \, \text{keV}$ per baryon. Typical values for $R_0$ are $\sim 650$ km, in reasonable agreement with the numerical studies of neutrino-driven winds [39, 41]. Since we are interested here in exploring the lower opacity case, we restrict our study to electron fractions above the threshold value for heavy r-process, $Y_{e,\text{thresh}} = 0.25$, see figure 8 in [34]. From the results shown in [41], we expect terminal wind velocities around $0.05c$. These numbers motivate our ‘wind ’ simulation. Once the parameters have been set, the density evolves according to $\rho(t) = \rho_0 (1 + v_w t/R_0)^{-3}$ and the temperature evolution is calculated using the HELMHOLTZ equation of state [101] according to the entropy change from nuclear reactions [7]. The explored wind parameters are given in table 2 and visualized in figure 4.

### 2.3. Nucleosynthesis calculations

For each simulation we perform a nucleosynthesis calculation to obtain the final abundance pattern and the time-dependent nuclear heating rate $\dot{\varepsilon}_{\text{nuc}}$ that is needed for the macronova models. We calculate an average out of 1000 randomly chosen hydrodynamic ejecta trajectories for each of our merger simulations (N1–N5). Since there is little variation between individual trajectories this is a fair representation of the overall ejecta dynamics. These average trajectories are used in the network calculations to obtain the nuclear energy generation rate $\dot{\varepsilon}_{\text{nuc}}$ that is used in the macronova calculation. For the dynamic ejecta in the NSBH cases (run B1–B3) and the winds we use the above described expansion model. For the NSBH cases we apply the parameters from the simulations of [26]. For the entropy and electron fraction they provide results for individual trajectories that are near $Y_e = 0.06$ and $s = 2 \, \text{keV}$ per baryon. We choose these values in our models, but stress that the nucleosynthesis (and therefore the macronovae) in this regime is insensitive to the exact numbers. The parameters for the dynamic ejecta are summarized in table 1. Since what we subsume under ‘winds’ can have different physical origins, see the discussion above, we vary the wind parameters in a wide range, see table 2.

Our baseline nucleosynthesis calculations are performed with a large nuclear reaction network WinNet [102, 103] that is based on the BasNet network [104]. It includes 5831 isotopes from nucleons up to $Z = 111$ between the neutron drip line and stability. The reaction rates

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9 To illustrate this, have a look at figure 8 in [34]: all trajectories with $Y_e < 0.15$ yield practically identical nucleosynthesis results.
are from the compilation of [105] for the finite range droplet model (FRDM) [106] and the weak interaction rates (electron/positron captures and $\beta$-decays) are from [107] and [108]. In addition, the distribution of fission fragments from [109] is used as a default. Finally, we adopt neutron capture and neutron-induced fission rates of [110] and $\beta$-delayed fission probabilities of [111].

The nucleosynthesis in the dynamic ejecta has been found to be extremely robust against variations of astrophysical parameters [34], but the final abundance patterns are still sensitive to variations of the nuclear physics such as fission distributions, $\beta$-decay rates and nuclear mass formulae [94, 112–116]. This is because, for example, different nuclear mass formulae that reproduce -where known- nuclear masses well differ from each other when extrapolated towards the neutron dripline where much of the r-process is happening. The variation of nuclear physics can also lead to different predictions of the nuclear heating rate during the macronova time scale and therefore impact on light curve predictions. In [115], this effect was investigated with four nuclear mass models for a particular trajectory from a NSNS merger simulation. It was found that the late-time nuclear heating rate can be influenced by the population

| Run | $m_W$ | $Y_e$ | $v_{W,\infty}$ | $E_{\text{kin}}$ | $X_{\text{lan}}$ | $X_{\text{act}}$ | Comment |
|-----|-------|-------|---------------|----------------|----------------|----------------|---------|
| Wind 1 | 0.01 | 0.30 | 0.05 | $2.2 \times 10^{49}$ | $1.61 \times 10^{-7}$ | $< 10^{-15}$ | Perego et al (2014) |
| Wind 2 | 0.01 | 0.25 | 0.05 | $2.2 \times 10^{49}$ | $6.30 \times 10^{-5}$ | $< 10^{-15}$ | Perego et al (2014) |
| Wind 3 | 0.01 | 0.35 | 0.05 | $2.2 \times 10^{49}$ | $< 10^{-15}$ | $< 10^{-15}$ | Perego et al (2014) |
| Wind 4 | 0.05 | 0.25 | 0.05 | $1.1 \times 10^{30}$ | $2.41 \times 10^{-4}$ | $< 10^{-15}$ | unb. torus |
| Wind 5 | 0.05 | 0.30 | 0.05 | $1.1 \times 10^{30}$ | $2.45 \times 10^{-7}$ | $< 10^{-15}$ | unb. torus; low-$\alpha$ |
| Wind 6 | 0.05 | 0.35 | 0.05 | $1.1 \times 10^{30}$ | $< 10^{-15}$ | $< 10^{-15}$ | unb. torus; low-$\alpha$ |
| Wind 7 | 0.05 | 0.25 | 0.10 | $4.5 \times 10^{50}$ | $4.28 \times 10^{-5}$ | $< 10^{-15}$ | |
| Wind 8 | 0.05 | 0.30 | 0.01 | $4.5 \times 10^{48}$ | $1.57 \times 10^{-4}$ | $< 10^{-15}$ | |
| Wind 9 | 0.10 | 0.25 | 0.10 | $9.0 \times 10^{50}$ | $7.70 \times 10^{-5}$ | $< 10^{-15}$ | |
| Wind 10 | 0.01 | 0.25 | 0.10 | $9.0 \times 10^{49}$ | $3.49 \times 10^{-5}$ | $< 10^{-15}$ | |
| Wind 11 | 0.01 | 0.25 | 0.25 | $5.9 \times 10^{50}$ | $2.13 \times 10^{-2}$ | $6.34 \times 10^{-6}$ | |
| Wind 12 | 0.01 | 0.25 | 0.50 | $2.8 \times 10^{51}$ | $7.50 \times 10^{-2}$ | $1.66 \times 10^{-5}$ | |
| Wind 13 | 0.10 | 0.35 | 0.01 | $8.9 \times 10^{48}$ | $< 10^{-15}$ | $< 10^{-15}$ | |
| Wind 14 | 0.10 | 0.30 | 0.05 | $2.3 \times 10^{50}$ | $1.35 \times 10^{-7}$ | $< 10^{-15}$ | |
| Wind 15 | 0.20 | 0.35 | 0.01 | $1.8 \times 10^{49}$ | $< 10^{-15}$ | $< 10^{-15}$ | |
| Wind 16 | 0.20 | 0.30 | 0.05 | $4.5 \times 10^{50}$ | $8.39 \times 10^{-8}$ | $< 10^{-15}$ | |
| Wind 17 | 0.20 | 0.25 | 0.10 | $1.8 \times 10^{51}$ | $1.27 \times 10^{-4}$ | $< 10^{-15}$ | |
| Wind 18 | 0.01 | 0.35 | 0.01 | $8.9 \times 10^{47}$ | $< 10^{-15}$ | $< 10^{-15}$ | |
| Wind 19 | 0.05 | 0.25 | 0.25 | $2.9 \times 10^{51}$ | $3.91 \times 10^{-4}$ | $< 10^{-15}$ | |
| Wind 20 | 0.10 | 0.25 | 0.25 | $5.8 \times 10^{51}$ | $4.95 \times 10^{-5}$ | $< 10^{-15}$ | |
| Wind 21 | 0.20 | 0.25 | 0.25 | $1.1 \times 10^{52}$ | $3.20 \times 10^{-5}$ | $< 10^{-15}$ | |
of trans-lead nuclei because of the additional heating via $\alpha$-decays. We explore here the two extreme cases out of the four mass formulae investigated in [115]—the FRDM mass formula which produces the smallest $\alpha$-decay heating rates and the 31-parameter Duflo–Zuker (DZ31) mass formula [117] which produces the largest, for a wider range of dynamical ejecta from NSNS and NSBH mergers listed in section 2.2. We note that while this represents the extreme cases of what has been explored so far, we cannot be entirely sure that this really brackets what may be realized in nature. This topic needs further exploration in the future.

The mass formula comparisons are performed with the network described in detail in [112]. It includes 7360 nuclei between the proton and neutron drip-lines up to charge number $Z = 110$. For the nuclear masses we use experimental values from the atomic mass evaluation [118] whenever available. Where available, we use experimental $\alpha\beta$-decay and spontaneous fission rates from the NUBASE [119] and NuDat2 databases (www.nndc.bnl.gov/nudat2/) and theoretical predictions otherwise. For the most relevant theoretical rates for the r-process, we employ the $\beta$-decay rates of [120], the neutron-capture and the reverse photo-dissociation rates of [112] for nuclei with $Z < 83$ for different mass models, the neutron-capture and neutron-induced fission rates of [110] for $Z > 83$ with FRDM mass model and Thomas–Fermi fission barriers, and the $\beta$-delayed and spontaneous fission rates of [121]. The fission fragment distributions are taken from [122] calculated by the ABLA code, which well reproduces available fission data and includes the possibility of neutron emission before and after fission.

### Macronovae

#### 2.4. Scaling relations

The luminosity of a MN rises as more material becomes visible and once all parts are transparent, one sees a decline in the luminosity dictated by the nuclear energy generation rate, $L \propto \epsilon_{\text{nuc}}m_\odot$. As a guidance for the following discussion, we provide...
the scaling laws that follow from simple arguments. After the neutron captures have ceased (~1 s after merger) the heating rate can to a reasonable approximation be described by a power law, $\dot{E}_{\text{rad}} \propto t^{-\alpha}$, with a power law index $\alpha \approx 1.3$ [12, 13, 34, 123]. The peak emission is reached when expansion and diffusion times are comparable which yields [17]

$$t_{\text{peak}} \approx 4.9d \left( \frac{\kappa}{10 \text{ cm}^2 \text{ g}^{-1}} \left( \frac{m_{\text{ej}}}{0.1 M_\odot} \right) \frac{0.1 c}{v_{\text{ej},\infty}} \right)^{1/2}$$

for the peak time, $\kappa$ being the effective opacity,

$$L_{\text{peak}} \approx 2.5 \times 10^{40} \text{ erg s}^{-1} \left( \frac{v_{\text{ej},\infty}}{0.1 c} \right)^{0.65} \left( \frac{m_{\text{ej}}}{0.01 M_\odot} \right)^{0.35}$$

for the bolometric luminosity and

$$T_{\text{eff}} \approx 2200 K \left( \frac{10 \text{ cm}^2 \text{ g}^{-1}}{\kappa} \right)^{0.4125} \left( \frac{0.1 c}{v_{\text{ej},\infty}} \right)^{0.0875} \left( \frac{0.01 M_\odot}{m_{\text{ej}}} \right)^{0.1615}$$

for the effective temperature where all exponents were evaluated with $\alpha = 1.3$.

### 2.4.2. Model 1

As our macronova model 1 (hereafter ‘MNmodel1’) we employ a model similar to the one used in [17] with a spherically symmetric homologously expanding radial density profile $\rho(\nu) \sim \rho_0 (1 - \nu^2 \nu_{\text{max}}^2)^3$. In this model, it is assumed that all macronova energy originates from the layer above the so-called diffusion surface, defined as the surface for which the average diffusion time is equal to the time from the merger. Any photons traveling below this diffusion surface are considered trapped and therefore invisible. The total luminosity is then taken to be the mass of this layer times the instantaneous radioactive nuclear heating rate; thermal emission from the layer is assumed to be lost to $PdV$ expansion work and therefore neglected. In this sense, our model provides a conservative lower bound for the macronova emission. The emitted spectrum has an effective temperature of the photosphere at optical depth $\tau_{\text{ph}} = 2/3$.

Practically, the models inherit from the hydrodynamic calculation the ejecta mass, the electron fraction and velocity and from the network calculation the instantaneous nuclear heating rate $\dot{E}_{\text{rad}}(t)$, assuming that a fixed fraction ($f_{\text{tot}} = 0.5 = \text{const}$) thermalizes at all times while the rest is lost. Note that in model 1 we always use the FRDM mass model [106].

A major uncertainty in the prediction of macronova transients are the opacities of the expanding r-process material, which unfortunately has the largest effect on the spectral energy distribution (SED), see equation (8). Based on atomic structure models, the authors of [96] argued that the opacity of expanding r-process material is dominated by bound-bound-transitions from those ions that have the most complex valence electron structure. They found in particular that even small amounts lanthanides have a large impact on the opacities. For example, the neodymium opacities exceed those of iron as long as their mass fraction $X_{\text{Nd}} > 10^{-4}$. From their studies based on four species, they conclude that a gray opacity of $\approx 10 \text{ cm}^2 \text{ g}^{-1}$ should be fairly effective for calculating bolometric light curves. Similar conclusions were reached in a study by [97]. However, [96] and, more recently, [98, 99] point out that even this may be an underestimate, and that the true value could even substantially larger.

We use the sum of lanthanide and actinide fraction, $X_{\text{lan}} + X_{\text{act}}$, to decide which opacity value to use. Whenever it exceeds a limiting value of $10^{-3}$, we use as fiducial value for the
opacity in the dynamic ejecta a value of \( \kappa = 10 \text{ cm}^2 \text{ g}^{-1} \), otherwise we use \( \kappa = 1 \text{ cm}^2 \text{ g}^{-1} \).

The dynamic ejecta show a very large lanthanide fraction of \( \approx X_{0.18} \text{lan} \) while this value varies widely for the different wind cases, see column seven in table 2, but is in almost all cases below the threshold value.

Since the current knowledge is based on expensive atomic structure calculations of so far only a few ions, the opacity value for strongly lanthanide-enriched material is likely subject to considerable uncertainties. We therefore also explore how our brightest case, N3 with DZ31-mass model, would appear in the case of a substantially larger opacity value (\( \kappa = 100 \text{ cm}^2 \text{ g}^{-1} \)).

2.4.3. Model 2. A major difference in macronova model 2 (‘MNmodel2’) is that we use time-dependent thermalization efficiencies. In addition, we can switch between the FRDM and the DZ31 nuclear mass model. All MNmodel2 calculations are performed with the [112] network. The effect of the DZ31 mass model is that a larger fraction of trans-lead nuclei are produced and, as shown below, this has a substantial impact on the nuclear heating rate at times around and after the macronova peak (\( t > 1 \text{ d} \)). The thermalization efficiencies have been explored in recent work [115, 123]. We apply here the thermalization efficiencies based on simple analytical estimates from the latter work. The thermalization efficiency for photons is estimated as

\[
f_\gamma(t) = 1 - \exp \left( \frac{1}{\eta_\gamma^2} \right),
\]

while for the massive particles (electrons, \( \alpha \)-particles and fission products) it reads

\[
f_k(t) = \frac{\ln(1 + 2\eta_k^2)}{2\eta_k^2},
\]

The quantity \( \eta \) is the ratio of time after the merger and the thermalization time scale of the considered particle, \( \eta_k = t/\eta_k \). We use for the different thermalization time scales

\[
\tau_\gamma = 1.40 m_5^{1/2} v_2^{-1} \text{ d}
\]

\[
\tau_e = 7.40 m_5^{1/2} v_2^{-3/2} \left( \frac{0.5 \text{ MeV}}{E_e} \right)^{1/2} \text{ d}
\]

\[
\tau_\alpha = 7.74 m_5^{1/2} v_2^{-3/2} \left( \frac{6.0 \text{ MeV}}{E_\alpha} \right)^{1/2} \text{ d}
\]

\[
\tau_{\text{fis}} = 16.77 m_5^{1/2} v_2^{-3/2} \left( \frac{125.0 \text{ MeV}}{E_{\text{fis}}} \right)^{1/2} \text{ d},
\]

where \( m_5 \equiv m_{\odot}(5 \times 10^{-3} M_\odot) \) and \( v_2 \equiv v_{\odot}/(0.2c) \). In the following, we evaluate the time scales \( \eta_k \) at the characteristic scaling energies given above. The total thermalization efficiency is then given by

\[
f_\text{tot}(t) = \frac{\epsilon_\beta(t) \left[ \zeta_\beta f_\beta(t) + \zeta_\alpha f_\alpha(t) + \epsilon_\delta(t) f_\delta(t) + \epsilon_{\text{fis}} t f_{\text{fis}}(t) \right]}{\epsilon_\beta(t) + \epsilon_\alpha(t) + \epsilon_\delta(t) + \epsilon_{\text{fis}}(t)},
\]

where \( \epsilon_\beta(t) \) and \( \epsilon_\alpha(t) \) are the heating rates of the \( \beta \)-decays and coherent p-emission, respectively, and \( \epsilon_\delta(t) \) is the heating rate from the decay of an identified neutron to a \( \alpha \)-particle. \( \epsilon_{\text{fis}}(t) \) is the heating rate from fission products.
where we use $\zeta = 0.45$ and $\zeta_e = 0.2$ [115]. The nuclear heating rate that enters the macronova calculation is then

$$\dot{e}_{\text{heat}}(t) = f_{\text{tot}} \left[ \dot{e}_{\beta}(t) + \dot{e}_{\alpha}(t) + \dot{e}_{\text{fis}}(t) \right].$$

(16)

3. Results

3.1. Nucleosynthesis

Unless mentioned otherwise, we refer to WinNet results with the FRDM mass model as a default. From our NSNS simulations we find an electron distribution reaching up to 0.3, but with the majority of matter being near 0.04. This results in a very robust r-process pattern up to and beyond the third, ‘platinum’ r-process peak near $A = 195$, see right panel in figure 5.

For the neutrino-wind models the picture is different: here the resulting abundance pattern is sensitive to the details and in particular to the electron fraction of a thermodynamic trajectory. Consequently, the resulting pattern varies strongly between different wind models. Due to their relative large electron fraction ($Y_e > 0.25$), none of them reaches substantially beyond nucleon numbers of $A = 130$.

An exhaustive comparison of nuclear network results is beyond the intention of this paper. Nevertheless, we perform a short comparison between the two networks [102, 112] for one case (dynamic ejecta of run N4, NSNS binary with 1.2 and 1.4 $M_\odot$; $t = 100$ d), simply in order to gauge by how much the nucleosynthesis results might be influenced by implementation details. In both cases we use the FRDM nuclear mass model. The final mass fractions are shown in figure 6, left panel. The overall pattern agrees reasonably well, but there are noticeable differences in the regime from $A \approx 90$ to 170, due to the treatment of fission. We also briefly explore which impact the nuclear mass model has on the final abundance pattern (run N4; Mendoza-Temis network; right panel of figure 6). A major difference is the substantially higher mass fraction of trans-lead ($A > 207$) nuclei when DZ31 is used. [115] found that $\alpha$-decays of these nuclei have a substantial impact on the late-time lightcurve, see below.

Overall, the DZ31 model shows a closer agreement with the solar r-process pattern, especially around the third r-process peak ($A \approx 195$).

We also briefly compare the total nuclear energy generation rate resulting from both networks (each time using FRDM, for all dynamic ejecta) in figure 7. There is good agreement between both networks over many orders of magnitude, only at very late stages we find for
some trajectories deviations of up to a factor of two. We leave more detailed comparisons of the effects of different input physics in the networks to future work.

3.2. Thermalization efficiency and nuclear heating rates

Since the various reaction products thermalize on different time scales, see equations (11)–(14), the nuclear mass model has also an impact on the total thermalization efficiencies. Since $\alpha$-decays (due to translead nuclei) are enhanced, we find noticeable differences in the overall efficiencies, see figure 8. In all the cases $f_{\text{tot}}$ remains approximately constant around $\sim 0.7$, but decreases substantially slower at late times (for $t > 3$ d) for the DZ31 mass model$^{10}$. Closely related, for DZ31 the net heating rate, equation (16), is significantly different, see figure 7, right panel: at late times ($t > 1$ d) it exceeds the FRDM-results by up to an order of magnitude. This has a serious impact on the resulting macronova lightcurves around the time of the peak emission.

3.3. Optical and near-IR macronova lightcurves

One of the major objectives of our work is to explore the detectability of macronovae using current and future wide field-of-view optical and near-IR facilities, either as a result of follow-up of LIGO triggers or through independent transient searches. Figures 9 and 10 show the expected optical and near-IR lightcurves in absolute magnitudes for the dynamical ejecta of our brightest NSNS (N5) and NSBH (B3) mergers, respectively. Each time, we show the results from our macronova model 1 (‘MNmodel1’; see section 2.4.2) and for model 2 (‘MNmodel2’; see section 2.4.3), once for the FRDM and once for the DZ31 nuclear mass model. Throughout this work we use LSST grizy filters and 2MASS JHK and magnitudes in the AB-system.

The dynamic ejecta in N5 has a mass ratio that deviates substantially from unity ($q = 0.78$), but is consistent with the currently known mass ratios of NSNS binaries. For example,

\[^{10}\] Our results here differ somewhat from [115] in the sense that the ejecta are denser/slower. Therefore, the efficiencies can be higher in first few days and can actually increase if the $\alpha$-decays do so.
the observed value of J0453 + 1559 is $q = 0.75$ [124] and the recently discovered PSR J1913 + 1102 could have an even lower mass ratio [125]. The left panel of figure 9 shows the results for the macronova model 1, the middle panel shows MNmodel2 with the FRDM mass formula and the right panel refers to MNmodel2 with DZ31. The general trends with a fainter and faster lightcurve in the bluer bands is apparent, while the near-infrared (NIR) lightcurves can stay bright for several weeks. We typically have $-11.5$ at peak in the g band versus $-13.8$ in the K band. The MNmodel2 results are about 0.7 magnitudes brighter at peak, but decay faster at later times. Both effects are mainly due to the time variation of the thermalization efficiency, see figure 8. As expected from the enhanced net heating rate at late times (see figure 7) the DZ31 mass model yields peak magnitudes that are another 0.8 magnitudes brighter than for the FRDM case. At the same time, both runs using MNmodel2 are significantly redder in the optical, being about one magnitude fainter at peak in the g band. Additionally, their g-band lightcurves peak much earlier—only half a day after the merger versus about three days for MNmodel1.

Figure 10 shows the predictions for run B3 (1.2 $M_{\odot}$ NS and a 7.0 $M_{\odot}$ BH with a dimensionless spin parameter $\chi = 0.9$). The BH mass of 7 $M_{\odot}$ is close to the expected peak of the BH mass distribution [126], but the spin is admittedly high. However, if we are interested in NSBH systems that are able to launch a short GRB, we need a large BH spin ($\chi \approx 0.9$) in the
first place in order to form an accretion torus. Otherwise the neutron star is essentially swallowed whole (see, for example, the discussion in section 5.3 of [29] and references therein) and no GRB can be launched. Moreover, measured BH spins in high-mass x-ray binary systems tend to have large spin values \( \chi > 0.85 \), and these systems are the likely progenitors of NSBH binaries [127].

In the most favorable NSBH case (MNmodel2, DZ31) a \( K \)-band peak magnitude brighter than \(-16\) is reached. The \( g \) band reaches a similar brightness as for the NSNS case, but declines on a faster time scale for MNmodel1. Again, the objects are quite red, with the brightest magnitudes and more long lived light curves in the NIR.

Figure 11 shows the lightcurves for selected wind models (MNmodel1). The parameters of the first shown model (wind3) are guided by the simulations of the neutrino-driven winds in the aftermath of a 1.4–1.4 \( M_\odot \) merger by [41], where the central remnant survives for at least a few hundred milli-seconds before collapsing into a BH. With only very few available wind simulations, we consider this for now as representative for neutrino-driven winds from a NSNS merger. We consider the parameters of the wind model shown in the second panel (wind4) as representative for unbound accretion disk material [40, 43, 94]. The last shown wind model (wind9) is a rather extreme case where 0.1 \( M_\odot \) is ejected at 0.1c, while still having a low opacity due to the large electron fraction \( Y_e = 0.25 \); the lanthanide fraction in this case is \( \sim 3.5 \times 10^{-5} \), see column six in table 2).
For an efficient comparison, we have summarised the basic properties of the lightcurves for all models considered in tables 3 and 4. In addition to the peak magnitudes in LSST $r$ and 2MASS $J$ band, we calculated the time for which the lightcurve is within 1 magnitude of peak in those filters. All the calculated light curves available at: http://snova.fysik.su.se/transient-rates/

3.4. Detection feasibility

Next, we estimate the expected number of observable macronovae by integrating the expected rate of neutron star mergers (NSMs) $R_{\text{NSM}}$ over the comoving volume in which the resulting macronova is observable. The expected number of observable macronovae then becomes:

$$n_{\text{MN}} = \int_{z < z_{\text{max}}} R_{\text{NSM}}(1 + z)^{-1} dV_c,$$

(17)

where $dV_c$ is the comoving volume element and $z_{\text{max}}$ is the maximum redshift at which the macronova is brighter than $m_{\text{lim}}$. These limiting magnitudes correspond to redshift limits $z_{\text{max}} < 0.08$ for the dynamic ejecta models and $z_{\text{max}} < 0.22$ for the wind models, for which only a few cases reach such a high redshift. Most wind models have $0.08 < z_{\text{max}} < 0.15$. At those distances the assumption that the volumetric rate of merger events is constant still holds at the required accuracy. Note, however, that the extrapolation to the right edge of figure 12 could be affected by a redshift dependence of the rate. For $R_{\text{NSM}}$ we adopted an ‘informed best guess’ value of $300 \text{ yr}^{-1} \text{ Gpc}^{-3}$, see figure 2. To study the detectability, the lightcurves in observer-frame grizJHK were calculated from the time-dependent SEDs of each run using the Python package sncosmo [128] to account for spectral redshift. Calculations of cosmological distances and volumes were performed using the Python package astropy [129]. We assumed a cosmology with $H_0 = 67 \text{ km s}^{-1} \text{ Mpc}^{-1}$ and $\Omega_m = 0.307$ [130].

The four panels in figure 12 show representative model predictions for the expected number of MNe per year as a function of limiting AB magnitude over the whole sky, normalised to a nominal volumetric rate $R_{\text{NSM}} = 300 \text{ yr}^{-1} \text{ Gpc}^{-3}$ (left axis) and the corresponding reach in distance (axis on the right). Dashed lines at 75 Mpc and 140 Mpc show the quoted sensitivity limit for GW detections for LIGO run [131, 132] for NSNS and NSBH respectively. We also indicate the limiting magnitudes of 60 and 180-second exposures for VISTA and LSST in $J$ and $K$ band and grizy, respectively. When determining the limiting magnitudes, we used the ESO Infrared Exposure Time Calculator\(^\text{11}\) for VISTA and a Python exposure time calculator.

\(^{11}\) www.eso.org/observing/etc/
for LSST$^{12}$, assuming a target signal-to-noise ratio of 5. We also indicate the expected limiting magnitudes for the Zwicky Transient Facility (ZTF) $gri \sim 21$.

Here we have concentrated on a selection of the available models (N2, N5, B3 with DZ31 and wind20). Given the assumed rate of NSNS mergers, we can deduce the number of observable events for any given survey. We have done the calculations for these two given existing/}

$^{12}$ https://github.com/lsst-sims/exposure-time-calc
upcoming facilities, but the community can use this for any of their favorite surveys. These calculations are also provided under http://snova.fysik.su.se/transient-rates/.

Figure 12 also allows to read off the relative efficiency to detect NSNS mergers in the different passbands. The lightcurves show that more flux is expected for longer wavelengths, indicating that near-IR surveys would be particularly suitable. However, given the difficulty for ground based instruments to observe at these wavelengths, it is reassuring that also optical telescopes have a good sensitivity for macronova detections. The figure shows that in this particular case, one minute exposures in the optical bands ($i$) can be as efficient in detecting the merger.

Figure 12 is based on the best currently available opacity information. The opacities for dynamic ejecta, however, have just begun to be explored and are therefore not very accurately known. We therefore explore also the hypothetical case that the dynamic ejecta opacities should be as large as $100 \text{ cm}^2 \text{ g}^{-1}$, see figure 13. If such extreme opacities should be realized in nature, this would substantially deteriorate the detection prospect for the dynamic ejecta.

The above feasibility study shows that there are indeed prospects for finding optical and NIR emission from macronovae with large telescopes. More than a handful of events would be expected each year, distributed over the entire sky. In fact, if a survey such as LSST would allow high enough cadence (about once per night) in multiple filters, one could even discover macronovae without a trigger. Such un-triggered searches have already started, but are hampered by both difficulties in subtracting faint transients on potentially bright host galaxies, and by the background of supernovae [133]. The most interesting prospects for finding them clearly comes in connection with triggers from either GW signals or high-energy emission (Gamma-ray bursts).

| run            | $r_{max}$ | $\Delta m_{1mag}$ ($r$)/$m_{max}$ | $\Delta m_{1mag}$ ($J$) | Comment |
|----------------|----------|-----------------------------------|--------------------------|---------|
| MNmodel2, FRDM |          |                                   |                          |         |
| ns12ns12 (N1)  | −11.28   | 1.72                              | −12.97                   | 3.71    |
| ns12ns14 (N4)  | −11.31   | 2.18                              | −13.15                   | 5.35    |
| ns13ns13 (N2)  | −11.24   | 2.07                              | −13.06                   | 4.87    |
| ns14ns14 (N3)  | −11.23   | 1.87                              | −12.93                   | 4.07    |
| ns14ns18 (N5)  | −11.27   | 3.39                              | −13.65                   | 6.85    |
| ns14bh7 (B1)   | −11.69   | 1.55                              | −13.69                   | 5.08    |
| ns14bh7 (B2)   | −11.61   | 1.77                              | −13.79                   | 6.67    |
| ns12bh7 (B3)   | −11.78   | 1.42                              | −14.08                   | 7.01    |
| MNmodel2, DZ31 |          |                                   |                          |         |
| ns12ns12 (N1)  | −11.65   | 3.30                              | −13.71                   | 6.03    |
| ns12ns14 (N4)  | −11.61   | 4.48                              | −13.95                   | 7.49    |
| ns13ns13 (N2)  | −11.67   | 4.59                              | −13.99                   | 7.31    |
| ns14ns14 (N3)  | −11.59   | 3.42                              | −13.64                   | 6.22    |
| ns14ns18 (N5)  | −11.77   | 6.63                              | −14.57                   | 9.60    |
| ns14bh7 (B1)   | −12.06   | 3.41                              | −14.53                   | 7.66    |
| ns14bh7 (B2)   | −12.04   | 5.60                              | −14.92                   | 9.64    |
| ns12bh7 (B3)   | −12.21   | 5.24                              | −15.29                   | 10.30   |
| ns12bh7 (B3)   | −8.32    | 0.22                              | −10.80                   | 7.76    |

Table 4. Same as table 3, but for the more sophisticated macronova model MNmodel2.
3.5. Follow-up of LIGO triggers

The optical follow-up of gravitational wave signals is currently a large effort within the transient community [134]. We can now relate our model predictions with ongoing surveys such as iPTF [135] and Pan-STARRS [136]. These are typically conducted in the $r$ band, reaching limiting magnitudes of $\sim 20 - 21$. This will be typical also for the upcoming very wide area search with ZTF. For the most recent GW signals, at a distance of 400 Mpc, the detectability is illustrated by figure 9 in [137].

To such large distances, none of our models for dynamical ejecta becomes bright enough to be detectable. Even the most optimistic wind models are below the detection threshold of [137] at these distances. However, note that the current LIGO configuration can detect NSNS mergers only out to 75 Mpc [131]. For such a distance, the wind model would be easily detected by the search such as that by [137], and also the more optimistic dynamical ejecta models (e.g. N5 MNmodel2 DZ31) could be detected by ZTF [135] if nearby enough. With the rate of such events assumed above, we would expect about one such event per year, and not all may be observable by optical telescopes.

The larger telescopes studied here (LSST/VISTA) will more easily detect any such counterpart to a GW trigger, and as the sensitivity of LIGO for NSNS mergers increases from 75 to 140 Mpc, the detectability is illustrated by figure 9 in [137].

**Figure 12.** Expected number of MNe for selected bright models for each of the three categories (for the dynamic ejecta we use the brighter results coming from the DZ31 mass formula). The markers show the expected depths for a 60 s (circle) and 180 s (diamond) exposures with VISTA ($J$ & $K$ band) or LSST ($grizy$). Square markers show the expected numbers for a depth of 21 mag in $grizy$ as expected for ZTF. On the right-hand side the $y$-axis shows the luminosity distance up to which the NS merger rate was integrated to obtain $n_{\text{NS}}$. The gray dashed lines show the LIGO range for discoveries of gravitational wave signals from NS-NS mergers (75 Mpc) and NS-BH mergers (140 Mpc). Note that the results are scaled to the ‘best guess’ for the NSNS merger rate of 300 yr$^{-1}$ Gpc$^{-3}$.

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The larger telescopes studied here (LSST/VISTA) will more easily detect any such counterpart to a GW trigger, and as the sensitivity of LIGO for NSNS mergers increases from 75 to 140 Mpc, the detectability is illustrated by figure 9 in [137].
200 Mpc, we will go from one event every second year to about 10 events per year, and these events would indeed be within range for these telescopes.

In this comparison we also note that the timescales for the brighter models is several days, and not the very steep declines from a fast neutron precursor envisioned by some groups, see figure 3 in [135] and our discussion in the introduction. This is of great importance in determining the observing strategy, and our models suggest that a nightly cadence would be suitable, with several nights needed to establish variability. On the other hand, pursuing the search for much more than 2 weeks, as in [137], seems not to be warranted by these models.

We thus conclude that forthcoming surveys with the 8-meter LSST (∼10 sq. deg field of view) and existing 4-meter VISTA (∼2 sq. deg) will be able to search for GW-EM counterparts for any expected trigger from LIGO involving at least one neutron star with relatively short exposures per field. With such surveys, search areas of hundreds of sq. deg may be searched in a single night.

The feasibility of detection of coincident events is therefore promising. For more nearby mergers within ∼30 Mpc, for most of the models considered also surveys like ZTF would have an excellent chance of capturing the lightcurves of macronovae in multiple filters.

3.6. GRB triggers for the macronova candidate

Another way to get a trigger for a macronova event is via high-energy emission. The best case for a macronova so far comes from [18]. The detection of the short GRB 130603B is currently the best macronova candidate, the GRB had a fast decaying optical afterglow and was well placed within the host galaxy [138] at redshift $z = 0.356$. The Hubble Space Telescope (HST) imaging of [18] detected an H-band candidate at $25.73 \pm 0.20$ (F160W) at 7 rest-frame days (9.5 in observer frame) past GRB. Simultaneous optical (F606W) observations did not detect the source down to 28.25 (95%).

Comparison to all models shows which ones are bright enough to explain this event as a macronova. In figure 15, upper panel, we display the magnitude in the mass-velocity plane.
The colors of the symbols indicate the magnitude difference with respect to the Tanvir et al. detection and thus show that several models are within a factor of 3 of the required luminosity (orange and red symbols). Therefore, if we interpret this event as a MN, we note that the opacities can not be very different from those used here. Even the otherwise most promising models (e.g. model B3) would not be able to account for these observations if the effective opacities should be as large as $\kappa = -100 \text{ cm}^2 \text{ g}^{-1}$.

The wind models provide sufficiently bright lightcurves for the macronovae. Most notably wind model 9 matches the $H$-band detection perfectly while model 7 is only $\sim 0.5$ mag too faint, which is not significant given the model uncertainties, see the right panel of figure 14. The parameters of these models are close to what is expected for unbound accretion torus material (see stage IV, VI and IX in figure 1). The $R$-band upper limit, on the other hand, constrains the wind used for model 8, for which 27.81 is predicted. Perhaps even more interesting is that the brighter dynamical ejecta models also provide light curves that match this event, see left and middle panels of figure 14.

4. Summary and discussion

The major aim of this study was to explore the detectability of radioactively powered electromagnetic transients in the aftermath of compact binary mergers. We have performed simulations of neutron star mergers to extract hydrodynamic trajectories of the ejected material. These have been complemented by trajectories representing neutron star black hole mergers and various forms of wind outflows that have been treated in a parametrized form. What we denote collectively as ‘winds’ can have different physical origins such as neutrino-driven outflows or the late-time release of accretion torus material.

Along these trajectories nuclear reaction network calculations were performed to extract the radioactive heating rates which, in turn, serve as input for the subsequent macronova modeling. We have compared two different reaction networks [102, 112] to ensure the robustness of our results. Since the r-process path for very neutron-rich ejecta meanders through a region of the nuclear chart where no experimental information is available such calculations are based on theoretical nuclear models. We have therefore explored the difference between two frequently used nuclear mass models, the finite range droplet model (FRDM) [106] and the 31-parameter model (DZ31) of [117]. For the aspects explored in this manuscript the main differences for the observability come from different predictions for the amount of $\alpha$-decaying trans-lead nuclei. We have further compared the results of two macronova models. The first one [17] uses the FRDM mass model and assumes a fixed thermalization efficiency of 50%. In our second, and more sophisticated model, we use time-dependent thermalization efficiencies.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure14.png}
\caption{Lightcurves for selected models transformed to the redshift of GRB 130603B ($z = 0.356$) in HST-band F160W (black) and F606W (red). The down-pointing triangles correspond to 95\% upper limits.}
\end{figure}
based on the recent work of [115] and we allow to switch between the FRDM and the DZ31 mass model. We have calculated the resulting macronova lightcurves in different bands for eight representative dynamic ejecta cases, see table 1, and for 21 wind models, see table 2.

We find that the nuclear mass model has a decisive impact on the results. In terms of nucleosynthesis, the DZ31 model reproduces the solar-system r-process abundances around the platinum peak ($A = 195$) with substantially higher accuracy than the FRDM model. It is worth emphasizing, however, that also other nuclear properties such as decay half-lives and fission yields have substantial uncertainties. With DZ31 one finds substantially larger amounts of trans-lead material and its $\alpha$-decay yields nuclear heating rates that are, at the times relevant for the macronova emission, about an order of magnitude larger than in the FRDM case, see figure 7, right panel. In addition, the thermalization efficiencies in the DZ31 cases are increased by factors of a few. As a result, the DZ31 macronova light curves are substantially brighter than those obtained with FRDM. It is worth noting that the two mass models represent the extreme cases of what has been explored so far (FRDM lowest; DZ31 highest), but we cannot be entirely sure that this really brackets the real nuclear heating rate. This topic needs more exploration in the future.

We find that the more promising dynamic ejecta models from NSNS mergers reach K-band peak magnitudes in excess of $-15$ (model N5 with a mass ratio close to the observed NSNS binary J0453 + 1559; see right panel of figure 9), while the brightest NSBH dynamic ejecta model reaches values beyond $-16$. The wind models with parameters inspired by neutrino-driven winds from an NSNS-merger peak below $-13$, while models mimicking unbound torus matter [43, 94] reach peak values of $-14$ and—with only slightly increased mass and velocity parameters—they can reach close to $-16$ (figure 11).

Since more flux is expected at longer wavelengths, near-IR surveys are particularly suitable for macronova detections. We note, however, that exposures in optical bands (i) may be as efficient in detecting such mergers. Since the transients are generally faint and other sources evolving on these time scales are abundant, the best detection chances come from events that are triggered by either GW-signals or GRB-emission. For our adopted ‘best guess merger rate’ of 300 Gpc$^{-3}$ yr$^{-1}$ larger telescopes such as LSST or VISTA should detect of order 10
events per year. For mergers within 30 Mpc ZTF would have an excellent chance of capturing macronovae in multiple filters.

We have also explored which of our models could be plausible explanations for the best-to-date macronova candidate, the transient observed in the aftermath of GRB 130603B [18, 19]. Adopting the more optimistic DZ31 models as our standard, several of the explored cases get close to this observation. Both dynamic ejecta models (the non-equal mass NSNS merger model N5 and the NSBH models B2 and B3) and several of our wind models (wind7, wind9, wind17, wind20 and wind21) produce transients with similar properties. We note, however, that none of our dynamic ejecta models would be able to account for this observation for the hypothetical case that the effective opacities would be substantially larger than our adopted value of 10 cm$^2$ g$^{-1}$ [96, 97, 139] or if the pessimistic heating rates of the FDRM mass model should be realized in Nature.

The results presented here are based on our currently best macronova models. These involve a chain of different calculations where each has its own challenges. For example, the calculations involve gravity, hydrodynamics, neutrino and nuclear reactions, thermalization efficiencies and an approximate treatment of radiation. For several of these ingredients there are substantial uncertainties, for example for the nuclear equation of the state, the nuclear physics near the neutron-dripline and the involved matter opacities. Moreover, we note that our current models are not using proper radiative transfer calculations. While the discussed results represent the current status of our models, they certainly can and will be improved in future.

Lightcurves, expected rates and redshift distributions are available under: http://snova.fysik.su.se/transient-rates/.

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