Dynamical ejecta synchrotron emission as a possible contributor to the changing behaviour of GRB170817A afterglow

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ABSTRACT
Over the past three years, the fading non-thermal emission from the GW170817 remained generally consistent with the afterglow powered by synchrotron radiation produced by the interaction of the structured jet with the ambient medium. Recent observations by Hajela \textit{et al.} 2021 indicate the change in temporal and spectral behaviour in the X-ray band. We show that the new observations are compatible with the emergence of a new component due to non-thermal emission from the fast tail of the dynamical ejecta of ab-initio binary neutron star (BNS) merger simulations. This provides a new avenue to constrain binary parameters. Specifically, we find that equal mass models with soft equation of state (EOS) and high mass ratio models with stiff EOS are disfavored as they typically predict afterglows that peak too early to explain the recent observations. Moderate stiffness and mass ratio models, instead, tend to be in good overall agreement with the data.

Key words: neutron star mergers – stars: neutron – equation of state – gravitational waves

1 INTRODUCTION

The GW170817 event marked the dawn of the era of multimessenger astronomy with compact binary mergers. This event was observed as gravitational-wave (GW) source, GW170817 (Abbott \textit{et al.} 2017a, 2019a,b); quasi-thermal electromagnetic (EM) transient, commonly referred to as kilonova, AT2017gfo (Arcavi \textit{et al.} 2017; Coulter \textit{et al.} 2017; Drout \textit{et al.} 2017; Evans \textit{et al.} 2017; Hallinan \textit{et al.} 2017; Kasliwal \textit{et al.} 2017; Nicholl \textit{et al.} 2017; Smartt \textit{et al.} 2017; Soares-Santos \textit{et al.} 2017; Tanvir \textit{et al.} 2017; Troja \textit{et al.} 2017; Mooley \textit{et al.} 2018; Ruan \textit{et al.} 2018; Lyman \textit{et al.} 2018); and short γ-ray burst (SGRB), GRB170817A (Savchenko \textit{et al.} 2017; Alexander \textit{et al.} 2017; Troja \textit{et al.} 2017; Abbott \textit{et al.} 2017b; Nynka \textit{et al.} 2018; Hajela \textit{et al.} 2019), detected by the space observatories Fermi (Ajello \textit{et al.} 2016) and INTEGRAL (Winkler \textit{et al.} 2011).

This SGRB was dimmer then any other events of its class. Different interpretations for its dimness and slow rising flux were proposed: off-axis jet, cocoon or structured jet. Now it is commonly accepted that GRB170817A was a structured jet observed off-axis (\textit{e.g.} Fong \textit{et al.} 2017; Troja \textit{et al.} 2017; Margutti \textit{et al.} 2018; Lamb & Kobayashi 2017; Lamb \textit{et al.} 2018; Ryan \textit{et al.} 2020; Alexander \textit{et al.} 2018; Mooley \textit{et al.} 2018; Ghirlanda \textit{et al.} 2019). The GRB170817A late emission, the afterglow, provided further information on the energetics of the event and on the properties of the circumburst medium (\textit{e.g.} Hajela \textit{et al.} 2019).

The non-thermal afterglow of GRB170817A has been observed for over three years, fading after its peak emission at \textit{\sim}160 days after merger. At the time of writing, 3.2 years past the merger, the post-jet-break afterglow is still being observed, albeit only in X-ray by Chandra (Hajela \textit{et al.} 2021) and in radio by VLA (Balasubramanian \textit{et al.} 2021), as its flux in optical wavelengths has decreased below the detection limit (Troja \textit{et al.} 2020). Up until 900 days after merger, the non-thermal emission in X-rays and radio have followed the typical post-jet-break afterglow decay, $t^{-\alpha}$ (Sari \textit{et al.} 1998). After 900 days, a flattening in the X-rays, \textit{i.e.}, behaviour divergent from the $t^{-\alpha}$ decay, was observed (Troja \textit{et al.} 2020). More recent observations by Chandra showed the emergence of a new rising component in X-ray, not accompanied by the increase in radio flux, indicating a change of the spectral behaviour of the afterglow (Hajela \textit{et al.} 2021; Balasubramanian \textit{et al.} 2021).

There are several possible explanations for this behaviour, that generally fall into two categories (Hajela \textit{et al.} 2021). The first one is related changes occurring in the same shock that produced the previously observed afterglow emission (Frail \textit{et al.} 2000; Piran 2004; Sironi & Giannios 2013; Granot \textit{et al.} 2018; Nakar 2020) and include the transition of the blast wave to the Newtonian regime, energy injection into the blast wave, change of the interstellar medium (ISM) density, emergence of a counter-jet emission, and evolution of the microphysical parameters of the shock. The second category tight to the emergence of a new emitting component (Nakar & Piran 2011; Piran \textit{et al.} 2013; Hotokezaka & Piran 2015; Radice \textit{et al.} 2018c; Hotokezaka \textit{et al.} 2018; Kathirgamaraju \textit{et al.} 2019; Desai

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et al. 2019; Nathani et al. 2021; Ishizaki et al. 2021, e.g.) and includes afterglow from the decelerating ejecta, produced at merger or/and after, and emission powered by accretion onto a newly formed compact object. Notably, while the fall-back accretion scenario provides a tentative explanation for the X-ray excess in the observed spectrum, it requires a suppression mechanism at earlier times, e.g., suppression of the fall-back due to $r$-process heating (e.g. Desai et al. 2019; Ishizaki et al. 2021), as the earlier emission from GW170817 was consistent with the structured off-axis jet afterglow (Troja et al. 2020; Hajela et al. 2019; Hajela et al. 2021, e.g.). The kilonova afterglow, on the other hand, is a more straightforward explanation, as the kilonova itself has been observed and the emergence of its afterglow is only natural. Additionally, the X-ray excess, or in other words, stepper electron spectrum with lower $p$, is expected for non-relativistic material at the forefront of the dynamical ejecta which can be explained by an emerging non-thermal afterglow emission from the neutron star (NS) EOS (Margalit & Metzger 2017; Bauswein 2020; Levkov et al. 2019; Bernuzzi et al. 2020; Nedora et al. 2021b). These simulations were performed with the GR hydrodynamics code WhiskyTHC (Radice & Rezzolla 2012; Radice et al. 2014a,b, 2015), and included neutrino emission and absorption using the M0 method described in Radice et al. (2016) and Radice et al. (2018c), and turbulent viscosity of magnetic origin via an effective subgrid scheme, as described in Radice (2017) and Radice (2020). The impact of viscosity on the dynamical ejecta properties was investigated in Radice et al. (2018b), Radice et al. (2018c), and Bernuzzi et al. (2020), while the importance of neutrinos for determining the ejecta properties was discussed in Radice et al. (2018c) and (Nedora et al. 2021b), where it was shown that neutrino reabsorption increases the ejecta mass and velocity – two main quantities for the kilonova afterglow. All simulations were performed using finite temperature, composition dependent nuclear EOSs. In particular, we employed the following set of EOSs to bracket the present uncertainties: DD2 (Typel et al. 2010; Hempel & Schaffner-Bielich 2010), BLh (Logoteta et al. 2021), LS220 (Lattimer & Swesty 1991), SLy4 (Douchin & Haensel 2001; Schneider et al. 2017), and SFHo (Steiner et al. 2013). Among them, DD2 is the stiffest (thus providing larger radii, larger tidal deformabilities and larger NS maximum masses), while SFHo and SLy4 are the softest.

In this work we show that the observed X-ray behavior can be explained by an emerging non-thermal afterglow emission from the fast tail of BNS dynamical ejecta. We consider state-of-the-art NR simulations targeted to the GW170817 event (i.e., with GW170817 binary chirp mass) and documented in Perego et al. (2019); Nedora et al. (2019); Bernuzzi (2020); Nedora et al. (2021b). We compute synthetic LCs from the simulated ejecta using semi-analytical methods and show that the peak time and flux are consistent with the recent observations from some of the models. This provides a new avenue to constrain the binary parameters, suggesting that the equal mass models with very soft EOS peak too early to be consistent with observed changing behaviour of the GRB170817A X-ray afterglow.

2 BINARY NEUTRON STAR MERGER DYNAMICS AND DYNAMICAL MASS EJECTION

NR simulations of BNS mergers provided a quantitative picture of the merger dynamics, mass ejection mechanisms and remnant evolution (e.g. Shibata & Hotokezaka 2019; Radice et al. 2020; Bernuzzi 2020).

We consider a large set of NR BNS merger simulations targeted to GW170817 (Perego et al. 2019; Endrizzi et al. 2020; Nedora et al. 2019; Bernuzzi et al. 2020; Nedora et al. 2021b). These simulations were performed with the GR hydrodynamics code WhiskyTHC (Radice & Rezzolla 2012; Radice et al. 2014a,b, 2015), and included neutrino emission and absorption using the M0 method described in Radice et al. (2016) and Radice et al. (2018c), and turbulent viscosity of magnetic origin via an effective subgrid scheme, as described in Radice (2017) and Radice (2020). The impact of viscosity on the dynamical ejecta properties was investigated in Radice et al. (2018b), Radice et al. (2018c), and Bernuzzi et al. (2020), while the importance of neutrinos for determining the ejecta properties was discussed in Radice et al. (2018c) and (Nedora et al. 2021b), where it was shown that neutrino reabsorption increases the ejecta mass and velocity – two main quantities for the kilonova afterglow. All simulations were performed using finite temperature, composition dependent nuclear EOSs. In particular, we employed the following set of EOSs to bracket the present uncertainties: DD2 (Typel et al. 2010; Hempel & Schaffner-Bielich 2010), BLh (Logoteta et al. 2021), LS220 (Lattimer & Swesty 1991), SLy4 (Douchin & Haensel 2001; Schneider et al. 2017), and SFHo (Steiner et al. 2013). Among them, DD2 is the stiffest (thus providing larger radii, larger tidal deformabilities and larger NS maximum masses), while SFHo and SLy4 are the softest.

When NSs collide and merger, matter is ejected through a number of different physical processes, gaining enough energy to become gravitationally unbound. In particular, the matter ejected within a few dynamical timescales (i.e., $\sim 10$ ms) after merger by tidal torques and hydrodynamics shocks driven by core bounces is called dynamical ejecta. It was found that, within the velocity distribution of the dynamical ejecta, some simulations contain also a very fast tail with $v_{\|} \geq 0.6c$ (Piran et al. 2013; Hotokezaka et al. 2013; Kyutoku et al. 2014; Metzger et al. 2015; Ishii et al. 2018; Hotokezaka et al. 2018; Radice et al. 2018c).

The extensive analysis of this tail and its origin in a sample of NR simulations showed that the total mass of this tail depends on the binary parameters and on the NS EOS, but it is typically $\sim 10^{-6} - 10^{-5} M_{\odot}$. This fast tail can be decomposed into two components: the early fast ejecta that are channeled to high latitudes and that originate at the collisional interface of, predominantly, equal mass models with soft EOSs; and the late fast ejecta, that are largely confined to the plane of the binary, and are driven by the shock breakout from the ejecta after the first core bounce (Radice et al. 2018c).

In the following we consider the fast ejecta tail in the set of GW170817 targeted simulations (see above). The ejecta properties are extracted from the simulations at a coordinate radius of $R = 300 G/c^2 M_{\odot} \approx 443$ km from the center corresponding to the furthest extraction radius available. This ensures that the ejecta had the longest possible evolution inside the computational domain. This is also consistent with Radice et al. (2018c). The simulations were performed at a standard resolution with a grid spacing at the most refined grid level $\Delta x \approx 178$ m. We also performed several simulations with higher resolution, $\Delta x \approx 123$ m, to assess the resolution effects on ejecta properties.
Ejection mechanism and properties of the fast tail of the ejecta shown for three simulations, with two EOSs: BLh and SFHo and two mass ratios: \( q = 1.00 \) and \( q = 1.22 \). The upper panel in each plot shows the time evolution of the maximum density in the simulation (green curves) and the mass flux of the ejecta with asymptotic velocities exceeding 0.6c (red curves). The bottom panel shows the mass histogram of the fast ejecta tail as a function of time. In both panels the outflow rate and histograms are computed at a radius of \( R = 443 \) km and shifted in time by \( t = (\nu_{\text{fast}})^{-1} \), \((\nu_{\text{fast}})\) being the mass averaged velocity of the fast tail at the radius \( R \). The plot shows that most of the fast ejecta are generally produced at first core bounce with a contribution from the second in models with soft EOSs.

Table 1. Properties of the fast tail of the dynamical ejecta (that has velocity \( \nu > 0.6c \)) for a list of NR simulations from Nedora et al. (2021b) for which this ejecta is found. Columns, from left to right, are: the EOS, reduced tidal parameter, mass ratio, mass of the fast tail, its mass-averaged electron fraction, and velocity, and the RMS half-opening angle around the binary plane. Asterisk next to EOS indicate a model with subgrid turbulence.

| EOS      | \( \Lambda \) | \( q \) | \( M_{\text{ej}} \) \([M_\odot]\) | \( \langle \nu_e \rangle \) \([c]\) | \( \langle \theta_{\text{RMS}} \rangle \) \([\text{deg}]\) |
|----------|----------------|--------|-------------------|-----------------|-----------------|
| BLh*     | 541            | 1.00   | 1.52 \times 10^{-6} | 0.25            | 0.63            | 59.55           |
| BLh      | 541            | 1.00   | 2.53 \times 10^{-5} | 0.32            | 0.68            | 30.05           |
| BLh      | 539            | 1.34   | 1.37 \times 10^{-6} | 0.28            | 0.62            | 40.73           |
| BLh*     | 539            | 1.34   | 8.02 \times 10^{-7} | 0.20            | 0.63            | 18.05           |
| BLh      | 540            | 1.43   | 1.19 \times 10^{-8} | 0.32            | 0.60            | 60.74           |
| BLh*     | 543            | 1.54   | 1.22 \times 10^{-6} | 0.31            | 0.62            | 61.23           |
| BLh      | 538            | 1.66   | 1.25 \times 10^{-6} | 0.32            | 0.62            | 51.40           |
| BLh*     | 532            | 1.82   | 6.40 \times 10^{-7} | 0.36            | 0.74            | 67.44           |
| DD2*     | 853            | 1.00   | 6.65 \times 10^{-6} | 0.28            | 0.63            | 12.80           |
| DD2      | 853            | 1.00   | 9.65 \times 10^{-7} | 0.28            | 0.63            | 23.27           |
| DD2*     | 847            | 1.20   | 4.19 \times 10^{-7} | 0.24            | 0.61            | 20.02           |
| DD2      | 846            | 1.22   | 1.34 \times 10^{-7} | 0.26            | 0.65            | 16.90           |
| LS220*   | 715            | 1.00   | 1.20 \times 10^{-7} | 0.36            | 0.61            | 47.76           |
| LS220    | 715            | 1.00   | 3.40 \times 10^{-8} | 0.26            | 0.61            | 70.38           |
| LS220    | 714            | 1.16   | 1.17 \times 10^{-8} | 0.28            | 0.62            | 43.90           |
| LS220*   | 714            | 1.16   | 4.33 \times 10^{-8} | 0.28            | 0.63            | 36.21           |
| LS220*   | 717            | 1.11   | 7.57 \times 10^{-7} | 0.35            | 0.66            | 43.25           |
| LS220    | 710            | 1.43   | 1.01 \times 10^{-5} | 0.37            | 0.66            | 76.98           |
| LS220    | 707            | 1.66   | 8.39 \times 10^{-7} | 0.38            | 0.73            | 52.94           |
| SFHo*    | 413            | 1.00   | 3.97 \times 10^{-6} | 0.31            | 0.67            | 52.76           |
| SFHo     | 413            | 1.00   | 2.92 \times 10^{-5} | 0.30            | 0.66            | 53.13           |
| SFHo*    | 412            | 1.13   | 8.40 \times 10^{-5} | 0.23            | 0.69            | 25.36           |
| SFHo     | 412            | 1.13   | 4.54 \times 10^{-5} | 0.29            | 0.67            | 31.69           |
| SLy4*    | 402            | 1.00   | 5.13 \times 10^{-5} | 0.31            | 0.67            | 48.85           |
| SLy4     | 402            | 1.00   | 3.21 \times 10^{-5} | 0.30            | 0.69            | 44.02           |
| SLy4*    | 402            | 1.13   | 1.70 \times 10^{-4} | 0.20            | 0.67            | 24.02           |

Notably, not all simulations are found to host a measurable amount of fast ejecta. Specifically, we find the absence of the fast ejecta component in simulations with stiff EOS and relatively high mass ratio, \( q = M_1/M_2 \geq 1 \), where \( M_1 \) and \( M_2 \) are the gravitational masses at infinity of the primary and secondary NSs respectively. The dynamical ejecta velocity distribution from these models shows a sharp cut-off at \( \leq 0.5c \). The absence of fast ejecta can be understood from the fact that at large \( q \) the ejecta are dominated by the tidal component, whose speed is largely set by the NSs velocities at the last orbit and the system escape velocity. Additionally, our models with large mass ratio (with fixed chirp mass) experience prompt collapse with no core bounce (Bernuzzi et al. 2020).

The production mechanism of the fast ejecta tail is shown in Fig. 1. Here, we define the fast ejecta tail to consist of material with asymptotic velocity \( \nu > 0.6c \), following Radice et al. (2018c). However, we remark that the choice of the velocity threshold 0.6c is mostly conventional. We also remark that the synchrotron light curves are computed using the full velocity structure of the ejecta, so this choice does not have any impact on our results. We find that in our sample of simulations the ejection of mass with velocity \( \nu > 0.6c \) coincides with core bounces, in agreement with previous findings by Radice et al. (2018c). In models with moderately soft EOS or large mass ratio, e.g., the equal mass BLh EOS model or the unequal mass models with softer EOS, e.g., SFHo EOS model, most of the ejecta originate at the first bounce. However, in equal mass models with very soft EOS, e.g., the equal mass SLy4 EOS model, we find that additional mass ejection occurs at the second bounce. Notably, while the first-bounce component is generally equatorial, the second-bounce component is more polar. This might be attributed to the increased baryon loading of the equatorial region resulting from the slow bulk of dynamical ejecta and with the disc forming matter.

The presence of the fast tail is robust and is not affected by resolution. The mass of the fast tail, \( M_{\text{ej}}(\nu > 0.6c) \), however, does have a resolution dependency, and we find that \( M_{\text{ej}}(\nu > 0.6c) \) changes by a factor of a few between simulations at standard and high resolutions. A larger sample of simulations performed at high resolutions is required to assess this uncertainty more quantitatively. The mean value of the fast tail mass is \( M_{\text{ej}}(\nu > 0.6c) = (2.36 \pm 3.89) \times 10^{-5} M_\odot \), where we also report the standard deviation. Other properties of the fast tail, such as velocity, electron fraction and angular
distribution, are more robust with respect to resolution, similarly to what is observed for the total dynamical ejecta (Nedora et al. 2021b).

We report the ejecta properties of simulations performed with standard resolution Tab. 1. We find that for most models, the mass averaged velocity of the fast tails, $\nu_\infty (v>0.6 \, c)$, is close to 0.6c with models with softer EOSs displaying higher velocities. The mass-averaged electron, $Y_e (v>0.6 \, c)$ is generally above 0.25, indicating that these ejecta were shock-heated and produced by neutrinos. High average electron fraction implies that weak r-process nucleosynthesis would occur, producing elements up to the 2nd r-process peak (Lippuner & Roberts 2015).

The total kinetic energy of the fast tail, $E_k (v>0.6 \, c)$, is shown in top panel of the Fig. 2. The error bars cover a conservative $\sim 1$ order of magnitude, that is obtained by considering the resolution dependency of the fast ejecta mass and velocity, and by assuming the same error measures adopted in Radice et al. (2018c). The figure shows that the total kinetic energy of the fast tail ranges between $\sim 10^{50}$ erg and $\geq 10^{50}$ erg. Overall, the kinetic energy of the fast tail does not show a strong dependency on the EOS, even if very soft EOSs (like SLy4 and SFHo) tend to have larger energies. The dependency on the mass ratio is more prominent, especially for the SLy4, SFHo and LS220 EOSs, where for the latter, the $E_k (v>0.6 \, c)$ rises by $\sim 3$ orders of magnitude between $q=1$ and $q=1.7$. Notably, for the BLh EOS models, the total kinetic energy does not change with the mass ratio.

In the lower panel of the Fig. 2 we show the RMS half-opening angle of the fast ejecta around the orbital plane. We assume a conservative error of 5 degrees, motivated by the comparison with higher resolution simulations. As the angular distribution of fast ejecta depends on the ejection mechanism, the figure allows to assess which mechanism dominates in each simulation. The fast ejecta tail is largely confined to the binary plane for the models with stiff EOS, e.g., DD2 EOS, where the core bounce ejection mechanism dominate. Meanwhile, in simulations with soft EOSs and high mass ratios, the fast ejecta has a more uniform angular distribution determined by an interplay between the core dynamics and finite temperature effects driving shocked outflow.

As the mass of the fast ejecta tail shows resolution dependency, so does its total kinetic energy. For three models for which the fast ejecta were found in both the standard and high resolution simulations, we find that $E_k (v>0.6 \, c)$ changes by at least factor of a few. The ejecta RMS half-opening angle of the orbital plane is less resolution dependent and its uncertainty is less than $\sim 50\%$.

Next we consider the distribution of the cumulative kinetic energy of the ejecta, defined as the kinetic energy of the ejecta whose mass is above a certain speed. We express it as a function of the $\Gamma \beta$ product, where $\beta$ is the ejecta velocity expressed in units of $c$, and $\Gamma = 1/\sqrt{1-\beta^2}$ is the Lorentz factor. We show $E_k (> \Gamma \beta)$
for representative set of models in Fig. 3. The plot displays that for most models the bulk of the kinetic energy is allocated to the low velocity matter, i.e., for $\Gamma \beta \leq 0.5$. Equal mass models show an extended high velocity tail, especially the $q = 1.00$ model with SLy4 EOS. The bottom panel of the Fig. 3 shows the cumulative kinetic energy distribution in terms of the $\beta^2$ product and angle from the plane of the binary for the $q = 1.00$ model with BLh EOS. The distribution is not uniform with respect to the polar angle. While the high energy tail extends up to the polar angle, the high velocity tail is more confined to the orbital plane. Notably, since the largest part (in mass) of the ejecta is equatorial it eludes the interaction with the $\gamma$-ray burst (GRB) collimated ejecta and expands into an unshocked ISM. The latter can decrease the ISM density and delay the peak of the synchrotron emission (Margalit & Piran 2020).

3 THE SYNCHROTRON EMISSION FROM EJECTA-ISM INTERACTION

Evaluating the synchrotron emission from the merger ejecta requires the calculation of the dynamical evolution of the blast wave as it propagates through the ISM. The dynamical evolution of a decelerating adiabatic blast wave can be described via the self-similar solutions. If the blast wave remains always relativistic, the Blandford-McKee (BM) solution (Blandford & McKee 1976) applies. If the blast wave remains always subrelativistic, the Sedov-Taylor (ST) solution (Sedov 1959) can be used. Another approach to compute the dynamics of the blast wave is to consider the hydrodynamical properties of the fluid behind the shock to be uniform within a given (thin) shell (e.g. Pe’er 2012; Nava et al. 2013). This thin homogeneous shell approximation allows to describe the entire evolution of the shell’s Lorentz factor from the free coasting phase (where the blast wave velocity remains constant) to the subrelativistic phase. However, there are limitations to this approach. Specifically, it was shown to differ from BM self-similar solution in the ultrarelativistic regime by a numerical factor (Panaitescu & Kumar 2000), and a self-similar solution of the non-relativistic deceleration (Huang et al. 1999). In application to the mildly relativist ejecta with velocity structure, the deviation was shown to be of order unity (Piran et al. 2013; Hotokezaka & Piran 2015).

We calculate the non-thermal radiation arising from the dynamical ejecta propagating into the cold ISM with the semi-analytic code PyBlastAfterglow. The method can be summarized as following. For a given distribution of energies as a function of velocity, we divide the ejecta into velocity shells and solve the adiabatic radial expansion of the ejecta in the thin shell approximation at each polar angle using the kinetic energy distributions discussed in Sec. 2. See also e.g., Piran et al. (2013) and Hotokezaka & Piran (2015) for similar treatments.

For the adiabatic evolution we adopt the blast wave dynamics formalism developed by Nava et al. (2013) where the evolution of the blast wave Lorentz factor is given by their equations 3-7, which we solve numerically via a 4th order adaptive step Runge–Kutta (RK) method. We neglect the effects of radiation losses and lateral spreading of the blast wave and focus on its evolution prior to and shortly after the onset of the deceleration. The EOS assumed is that of the ideal transrelativistic fluid, where the adiabatic index is given as a function of the normalized temperature (equation 11 in Pe’er 2012) which is computed adopting the polynomial fit (equation 5 in Service 1986). Next, we consider the forward shock propagating into the upstream medium as the blast wave expands. The bulk of the energy is being deposited into the non-thermal protons. Part of this energy is transferred to relativistic electrons via complex shock interactions. It is however possible to consider a simplified prescription for the transfer of energy from protons to electrons (e.g. Dermer & Chiang 1998). A fraction $\varepsilon_e$ and $\varepsilon_B$ of shock internal energy is assumed to be deposited into the relativistic electrons and magnetic field respectively. The injected electrons are assumed to have a power-law distribution $dN/d\gamma_e \propto \gamma_e^{-p}$, where $\gamma_e$ is the electron Lorentz factor, $p$ is the spectral index, a free parameter. The critical Lorentz factors of the spectrum are the minimum one, $\gamma_{\text{min}}$, and the critical one, $\gamma_c$, computed via standard expressions (equations A3 and A4 in Johannesson et al. 2006, respectively). Depending on the ordering of the $\gamma_{\text{min}}$ and $\gamma_c$, two regimes are considered, namely the fast cooling regime if $\gamma_{\text{min}} > \gamma_c$, and slow cooling regime otherwise (Sari et al. 1998).

The coming synchrotron spectral energy distribution (SED) is approximated with a smooth broken power law according to Johannesson et al. (2006), and computed with their equations A1 and A7 for the slow cooling regime and their equations A2 and A6 for the fast cooling regime. The characteristic frequencies are obtained from the characteristic Lorentz factors $\gamma_{\text{min}}$ and $\gamma_c$, via their equation A5.

The synchrotron self-absorption is included via flux attenuation (e.g. Dermer & Menon 2009). However, for the applications discussed in this paper, the self-absorption is not relevant as the ejecta remains optically thin for the emission $\geq 3$ GHz (e.g. Piran et al. 2013).

We compute the observed flux, integrating over the equal time arrival surface (EATS), following Lamb et al. (2018). For each segment of the blast wave, the time for the observer is evaluated via their equation 3, and then the observed, Doppler-shifted flux is obtained via their equation 2. See Salmonson (2003) for the detailed discussion of the method and Hajela et al. (2021) for a similar implementation.

In the ultrarelativistic regime, evolving a single velocity shell, the code was found to be consistent with afterglowpy (Ryan et al. 2020), while in the subrelativistic regime, modeling the kilonova afterglow, the code produces LCs consistent with the model of Hotokezaka & Piran (2015), which was applied to the BNS ejecta in Radice et al. (2018c).

It is however important to note that the methods discussed above for both the blast wave dynamics and the synchrotron emission become increasingly inaccurate as the blast wave decelerates and spreads, and as most of the electrons become subrelativistic. So we do not discuss the late-time emission after the LC peak. We discuss
different physics implemented in PyBlastAfterglow with application to both structured GRBs and analytic ejecta profiles elsewhere.

The free parameters of the model are chosen as follows. We assume the ISM density to be uniform within the range $n_{\text{ISM}} \in (10^{-3}, 10^{-2}) \, \text{cm}^{-3}$ (Hajela et al. 2019). The observational angle, defined as the angle between the line of sight and the polar axis of the BNS system, is $\theta_{\text{obs}} = 30$ deg (Abbott et al. 2017a). For the luminosity distance of NGC 4993, the host galaxy of GW170817, we adopt $41.3 \times 10^6$ pc with the redshift $z = 0.0099$ (Hjorth et al. 2017).

Recent Chandra observations showed that the emerging component in GRB170817A afterglow is accompanied by the onset of the spectral evolution (Hajela et al. 2021). The observation data analysis suggests that a lower value for the electron power law distribution slope is more favorable, $p = 2.05$, but at low significance level. Due to this uncertainties, we consider the following parameter ranges $\varepsilon_e \in (0.1, 0.2), \varepsilon_B \in (10^{-3}, 10^{-2}), p \in [2.05, 2.15]$. Additionally, the effect of a lower $p = 2.05$ is discussed in Hajela et al. (2021).

### 4 RESULTS

We show X-ray and radio LCs from several representative models in Fig. 4, alongside the latest GRB170817A observational data.

The LC shape is determined by the ejecta velocity and angular distribution. For instance, models with broad velocity distribution, such as equal mass model with SLy4 EOS (see Fig. 3), have a wide LC. This LC starts to rise very early (in comparison with the high mass ratio model) as the fast velocity shells decelerate and peak on a shorter timescales. However, models with with narrower velocity distribution, such as the model with LS220 EOS and $q = 1.43$, have a narrower LC that rises later ($\sim 10^3$ days after merger).

Generally, within the uncertain microphysics and ISM density, the kilonova afterglow from most models is in a good agreement with the new observations by Chandra. In particular, models with $1.00 < q < 1.82$ and moderately stiff EOSs show rise in flux $\geq 10^3$ days postmerger, in agreement with observations.

Fixing the microphysical parameters and ISM density to $n_{\text{ISM}} = 5 \times 10^{-3}$ g cm$^{-3}$, $\varepsilon_e = 0.1$ and $\varepsilon_B = 5 \times 10^{-3}$, we find that the flux at the LC peak, $F_{\nu,p}$, is the largest for soft EOSs such as SFHo and SLy4. Overall, however, the peak flux seems largely independent of EOS. The $F_{\nu,p}$ dependency on the mass ratio for soft EOSs is that, the higher the mass ratio, the lower the peak flux. This can be understood from the following considerations. While the ejecta total kinetic energy budget of these models increases with the mass ratio, the mass-averaged velocity decrease (see Fig. 5 in Nedora et al. (2021b)). And slower, more massive ejecta have lower peak flux. However, for stiffer EOSs, such as DD2 and LS220, the $q$ dependency is less clear.

The LC shape and the peak time depend weakly on the uncertain microphysics and $n_{\text{ISM}}$. Specifically, within $n_{\text{ISM}} \in (10^{-3}, 10^{-2}) \, \text{cm}^{-3}$, $t_p$ varies by a factor of a few. An additional source of uncertainties is the ejecta properties dependency on numerical resolution. However we estimate its effect on the LCs to be smaller than that of the unconstrained shock microphysics and $n_{\text{ISM}}$. Specifically, the $t_p$ changes by a factor of $\leq 2$, and $F_{\nu,p}$ changes within a factor of $\leq 4$.

### 5 DISCUSSION

In this work we analyzed a large set of NR BNS simulations performed with state-of-the-art NR code WhiskyTHC and targeted to GW170817. Simulations included the effects of neutrino emission and reabsorption, effective viscosity via subgrid turbulence, and microphysical finite-temperature EOSs.

We found that most simulations’ ejecta contains material with velocities $\gtrsim 0.6$ c, whose properties are in agreement with previous studies (Radice et al. 2018). However, in binaries with large mass ratio and/or prompt BH formation (that experience weaker/no core bounce at merger), the fast ejecta could be absent.

The latest GRB170817A observations by Chandra at $10^3$ days showed a changing afterglow behaviour in X-ray band (Hajela et al. 2021). We suggest that this change can be attributed to the emergence of the kilonova afterglow. In particular, we evolved the ejecta from NR simulations with a semi-analytic code and computed its synchrotron emission. We found that the synthetic LCs are in agreement with the emerging new component in the GRB170817A af-
tergnow within the range of credibility of the microphysical parameters and of the ISM density, $n_{\text{ISM}}$. Additionally, the change in GRB170817A afterglow has the following implications: the kilonova afterglow peak should be (i) later and (ii) brighter than the latest observations. The condition (ii) is not particularly strong, as the LC peak flux depends sensibly on the uncertain microphysical parameters. The (i) condition is more robust from that point of view the LC peak flux depends sensibly on the uncertain microphysical parameters and allows to gauge which NR models from our sample have afterglow predictions more supported by observations.

In Fig. 5 we show the time of the LC peak for all models (for fixed microphysics and $n_{\text{ISM}}$, see Tab. 3), including those with absent fast tail, as these models still have sufficient amount of mildly relativistic material to produce a bright afterglow. The time of the synchrotron LC peak, $t_p$, of all the models is distributed around $\sim 10^3$ days postmerger (except tidal disruption cases, as the model with $q = 1.82$ and BLh EOS). Models with small mass ratio tend to have $t_p < 10^3$ days, while more symmetric models point towards $t_p > 10^3$ days. This trend is more apparent for models with soft EOSs. The reason for this lies in the mechanism responsible for the the fast ejecta tail (see Sec. 2). As the mass ratio increases, the amount of the shocked ejecta component decreases, and so does the kinetic energy of the ejecta fast tail. The afterglow of the slower, more massive ejecta peaks later (e.g. Hotokuzeaka & Piran 2015). Indeed, the time of the LC peak depends primarily on the ejecta dynamics, the so-called deceleration time (e.g. Piran et al. 2013). Additionally, in Fig. 5 we show the lower limit on $t_p$ (the time of the latest observation). We observe that LCs of models with moderate amount of fast ejecta, e.g., asymmetric models with EOSs of mild stiffness, lie above the limit, while models with highly energetic fast tails, such as equal mass models with very stiff EOSs, peak too early. It would be interesting to combine these constraints with those obtained from the modeling of the thermal component of the kilonova. However, this is not presently possible because the thermal emission is expected to be dominated by slowly expanding secular winds from the merger remnant, which cannot be presently simulated in full-NR (see e.g. Nedora et al. 2019).

The main conclusion of our work is that the observed GRB170817A changing behaviour at $10^3$ days after merger can be explained by the kilonova afterglow produced by ejecta in ab-initio NR BNS simulations targeted to GW170817. Specifically, models that produce a mild amount of fast ejecta, those with moderately to large mass ratio and moderately stiff EOSs are favoured. The dominant uncertainties in our analysis are the ill-constrained microphysical parameters of the shock. Additionally, the systematic effects due to the finite resolution, neutrino treatment and EOSs might be important. A larger set of observations, that allows for a better assessment of shock microphysics, and a larger sample of high resolution NR simulations are required to investigate these uncertainties further. We leave this to future works.

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**Data Availability:** the lightcurves and table 1 are available at (Nedora et al. 2021a).

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![Figure 5. Peak time, $t_p$, for LC for all considered NR simulations. Dashed black line corresponds to the last observation of GRB170817A afterglow, where the rising flux implies that it is a lower limit on the kilonova afterglow. The microphysical parameters and ISM density for all models are fixed and given in the Tab. 3. The plot shows that in general the $t_p$ increases with mass ration and with softness of the EOS, except for the softest, DD2 EOS.](image-url)
