Jet–Cocoon Outflows from Neutron Star Mergers: Structure, Light Curves, and Fundamental Physics

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Abstract

The discovery of GW170817, the merger of a binary neutron star (NS) triggered by a gravitational wave detection by LIGO and Virgo, has opened a new window of exploration in the physics of NSs and their cosmological role. Among the important quantities to measure are the mass and velocity of the ejecta produced by the tidally disrupted NSs and the delay—if any—between the merger and the launching of a relativistic jet. These encode information on the equation of state of the NS, the nature of the merger remnant, and the jet launching mechanism, as well as yielding an estimate of the mass available for r-process nucleosynthesis. Here we derive analytic estimates for the structure of jets expanding in environments with different density, velocity, and radial extent. We compute the jet–cocoon structure and the properties of the broadband afterglow emission as a function of the ejecta mass, velocity, and time delay between merger and launch of the jet. We show that modeling of the afterglow light curve can constrain the ejecta properties and, in turn, the physics of neutron density matter. Our results increase the interpretative power of electromagnetic observations by allowing for a direct connection with the merger physics.

Key words: gamma-ray burst: general – gravitational waves – stars: neutron

1. Introduction

The discovery of GRB170817A in association with the gravitational wave (GW) source GW170817, produced by the merger of a double neutron star (NS; Abbott et al. 2017), has opened a new era in the study of short gamma-ray bursts (SGRBs) as well as of NSs, both separately, and, especially, when the information from GWs and the electromagnetic counterparts are combined. Prior to the association with this GW event, only the brightest SGRBs were followed at longer wavelengths. However, even under these conditions, the afterglows were generally dim, and full broadband coverage was rarely possible (Fong et al. 2015).

GRB170817, despite having an isotropic luminosity \( \sim 10^4 \) times lower than that of “standard” SGRBs \((L \sim 10^{51} \text{ erg s}^{-1})\), was very bright due to its relatively small distance of 40 Mpc, having the initial trigger in GWs rather than in \( \gamma \)-rays as for the standard cosmological SGRBs (Goldstein et al. 2017; Savchenko et al. 2017). A massive broadband follow-up campaign to GW170817 allowed for unprecedented data coverage, revealing a kilonova/macronova (e.g., Arcavi et al. 2017; Coulter et al. 2017; Cowperthwaite et al. 2017; Smartt et al. 2017; Soares-Santos et al. 2017), as well as the signatures of a relativistic jet (Alexander et al. 2017, 2018; Margutti et al. 2017, 2018; Lazzati et al. 2018; Mooley et al. 2018; Ghirlanda et al. 2019). The mass of the ejecta was inferred to be \( \approx 0.065 M_\odot \) by Kasen et al. (2017) and between 0.042 and 0.077 \( M_\odot \) by Perego et al. (2017) from modeling of the kilonova light curve. The former analysis assumed two components, while the latter assumed three components, each having different density, velocity, and opening angle. As pointed out by Barnes et al. (2016), key uncertainties in kilonova modeling include the emission profiles of the radioactive decay products—nonthermal \( \beta \) particles, \( \alpha \) particles, fission fragments, and \( \gamma \) rays—and the efficiency with which their kinetic energy is absorbed by the ejecta.

The ejecta mass is an important quantity to measure in BNS mergers: it contains information on the equation of state (EoS) of the NS (e.g., Hotokezaka et al. 2013), and is a favorable site for the production of the heaviest elements in the universe via the r-process (Lattimer & Schramm 1974). As LIGO and Virgo have begun a new observing run at higher sensitivity, more NS–NS events will be expected, and hence any additional diagnostics of the ejecta mass will be extremely valuable.

The role of the ejecta in determining the properties of the jet has been studied under different angles by a number of investigators, especially within the context of explaining the multiwavelength observations of GRB170817A (e.g., Murguia-Berthier et al. 2014; Nagakura et al. 2014; Lazzati et al. 2017b, 2018; Bromberg et al. 2018; Gottlieb et al. 2018; Lamb & Kobayashi 2018; Troja et al. 2018; Wu & MacFadyen 2018; Xie et al. 2018; Beniamini et al. 2019; Beniamini & Nakar 2019; Gottlieb et al. 2019; Geng et al. 2019; Kathergamaraju et al. 2019). These data were found to be incompatible with a top-hat jet, but rather required a “structured jet,” as produced by the interaction of a relativistic outflow with the ejecta generated during the merger process.

The role of the ejecta was also recently discussed (Barbieri et al. 2019) in the context of NS–BH (black hole) mergers, together with the role of the mass and spin of the BH, in determining the observable electromagnetic signatures accompanying the coalescence. Additionally, the interaction of the ejecta themselves with the surrounding interstellar medium is expected to produce a long-lasting radio flare (Hotokezaka & Piran 2015).

Given the importance of the dynamical ejecta in molding the properties of the relativistic outflow that propagates in it, in this paper we perform a detailed study aimed at connecting the
properties of this matter (and in particular its mass and velocity), with the jet–cocoon structure that develops in it, and the resulting afterglow emission in representative bands. The sensitivity of this radiation to the ejecta properties (as will be shown here) will provide an additional element linking observable quantities to fundamental physics.

Our paper is organized as follows. We begin (Section 2) by deriving analytic expressions for the structured jet properties, connecting the jet morphology to the properties of the merger ejecta in which it propagates. Aided by the results of hydrodynamic (HD) numerical simulations, we derive energy and Lorentz factor profiles as a function of the off-axis angle. We feed these profiles to an external shock synchrotron code and compute the corresponding afterglow light curves (Section 3). Finally, we discuss the connection between these observable quantities and the NS EoS (Section 4). We summarize and conclude in Section 5.

2. The Structured Jet Model

In this section we discuss an analytic framework to compute the properties of the structured jet from a binary NS merger, within the framework of the Kompaneets approximation (e.g., Begelman & Cioffi 1989; Matzner 2003; Lazzati & Begelman 2005; Bromberg et al. 2011, 2014; Gill et al. 2019). We first discuss the input of HD simulations in setting up an analytic profile for the energy and Lorentz factor as a function of the off-axis angle. Subsequently, we derive the scale angles and energy of the jet and cocoon components.

(a) The energy and Lorentz factor profiles. To describe the angular distribution of the relativistic outflow at large radii (well after the jet interaction with the nonrelativistic (NR) ejecta) we adopt an exponential model, as supported by the results of numerical simulations (Figure 1):

\[
\frac{dE}{d\Omega} = Ae^{-\theta/\theta_i} + Be^{-\theta/\theta_c},
\]

The constants \(A\) and \(B\) are determined by the conditions that \(A \int_{\Omega} e^{-\theta/\theta_i} d\Omega\) is the total energy in the relativistic jet and \(B \int_{\Omega} e^{-\theta/\theta_c} d\Omega\) is the total energy in the cocoon. The variables \(\theta_i\) and \(\theta_c\) indicate the scale angles for the jet and cocoon, respectively. Note that \(\theta_i\) is in general smaller than the injected jet opening angle \(\theta_{inj}\) due to HD recollimation.

A similar profile is adopted for the Lorentz factor, but the value is not allowed to be less than unity:

\[
\Gamma = \Gamma_{core}e^{-\theta/\theta_i} + \Gamma_c e^{-\theta/\theta_c} + 1,
\]

where \(\Gamma_{core}\) is a free parameter of the model. The cocoon maximum Lorentz factor \(\Gamma_c\) is instead the result of the complex interaction between the jet and the ejecta, and is mainly controlled by the mixing between the two components. This cannot be modeled analytically and we assume here that \(\Gamma_c = \Gamma_{core}/10\), a value that fits our numerical results well (see the bottom panel of Figure 1). Given the fact that we only have one simulation to compare to, this value should be considered tentative. It should be remembered, however, that the initial Lorentz factor does not significantly affect the afterglow light curve at times longer than the deceleration time.

We now focus on calculating the energy profile (top panel of Figure 1). To evaluate the values of the four structure constants \(A, B, \theta_i,\) and \(\theta_c\) we model the propagation of a jet through the NR ejecta. We adopt the Kompaneets approximation (Begelman & Cioffi 1989; Matzner 2003; Lazzati & Begelman 2005; Bromberg et al. 2011, 2014) and identify the scale angles \(\theta_i\) and \(\theta_c\) as the opening angles of the jet and cocoon at break-out.

(b) The NS ejecta. The ejecta structure described above is created by the interaction of a relativistic jet with NR ejecta from the merger. We model the jet as an NR wind with constant velocity \(v_w\) and mass-loss rate \(\dot{m}_w\) that is released during the merger and by the remnant a time \(\delta t\) before the jet is launched. The ejecta are assumed to be spherically symmetric within a solid angle \(\Omega_w\). The density of the ejecta is therefore given by:

\[
\rho_w = \frac{\dot{m}_w}{\Omega_w v_w^2}.
\]

(c) The jet. At time \(\delta t\) after the NR ejecta are launched, a jet is launched from the central engine. We consider a top-hat jet of luminosity \(L_j\) with uniform injection Lorentz factor \(\Gamma_{j,inj}\) within an injection opening angle \(\theta_{j,inj}\). The jet is injected very simply; its “structure” is acquired during the interaction with the ejecta.

One can qualitatively understand the jet propagation as follows. The jet head propagates through the ejecta driving a bow shock that inflates a cocoon of hot material that enshrouds the jet. The cocoon pressure acts on the jet, collimating it into a narrower angle. The cocoon is trapped as well inside the ejecta and drives a shock into them. The model has three fundamental unknowns: the velocity of the head of the jet \(v_{j,head}\), the opening angle of the jet at breakout \(\theta_j\), and the opening angle of the cocoon at breakout \(\theta_c\). The system can be solved because it is constrained by three pressure balances as will be discussed in part (d) below.

The jet propagates inside the ejecta and eventually breaks out when the outer radius of the ejecta coincides with the position

![Figure 1. Comparison among the energy and Lorentz factor profile from a numerical HD simulation (Lazzati et al. 2017), the double exponential best-fit function that inspired our model (see Equations (1) and (2)) and the profiles predicted by our model.](image)
of the jet head, i.e., when
\[ v_w(t_{\text{bo}} + \delta t) = v_{jh}t_{\text{bo}}. \] (4)
This yields a breakout time (measured from the launching of the jet)
\[ t_{\text{bo}} = \frac{v_w}{v_{jh} - v_w} \delta t, \] (5)
and a breakout radius (Murguia-Berthier et al. 2014)
\[ r_{\text{bo}} = v_{jh}t_{\text{bo}} = \frac{v_wv_{jh}}{v_{jh} - v_w} \delta t. \] (6)

(d) The pressure balances. The structures of the jet and cocoon are defined by three pressure balances: at the head of the jet, the ram pressure of the jet on the contact discontinuity \( (p_{\text{ram,j}}) \) is balanced by the ram pressure of the merger ejecta on the same contact discontinuity \( (p_{\text{ram,wc}}) \). At the jet–cocoon transition, the cocoon thermal pressure \( (p_c) \) is balanced by the jet pressure \( (p_j) \), by either the jet internal pressure or the ram pressure due to the deflection of the jet material. Finally, at the cocoon–ejecta transition, the thermal pressure of the cocoon is balanced by the ram pressure of the merger ejecta \( (p_{\text{ram,wc}}) \).

These conditions yield the equations:
\[ p_{\text{ram,j}} = p_{\text{ram,wc}}, \] (7)
\[ p_j = p_c, \] (8)
\[ p_c = p_{\text{ram,wc}}. \] (9)

This set of equation allows us to solve for the three unknowns, \( v_{jh}, \theta_j \), and \( \theta_c \).

The first pressure balance (Equation (7)) has been studied in detail previously (e.g., Martí et al. 1994; Matzner 2003; Lazzati & Begelman 2005; Bromberg et al. 2011) and yields the speed with which the jet head advances in the merger ejecta. In static ejecta, the velocity of the jet head reads
\[ v_{jh} = \frac{c}{1 + \sqrt{\frac{\rho_j}{\rho_w}}}, \] (10)
where we have made the assumption that the jet material moves with \( v_j \approx c \) and that the enthalpy density is dominated by the pressure component.

If the jet propagates in a moving medium, the velocity in Equation (10) needs to be relativistically added to the wind velocity to yield the jet head velocity (Murguia-Berthier et al. 2017; Matsumoto & Kimura 2018; Gill et al. 2019)
\[ v_{jh} = \frac{v_w + v_{j'}^c}{1 + \frac{v_wv_{j'}^c}{c^2}}. \] (11)

Equation (10) can be written in terms of unknown quantities and model parameters. The wind density is given by Equation (3) and, using \( 4p_j^\frac{1}{\gamma} r^2 = L_j/(\Omega_c c) \), we obtain:
\[ v_{j'}^c = \frac{c}{1 + \frac{\omega_c}{\nu_j^r} \Omega_c r^3}. \] (12)

For the second pressure balance, we assume that the cocoon energy is dominated by radiation so that \( p_c = E_c/(3V_c) \), where \( V_c = \Omega_c r^3 \) is the cocoon volume and
\[ E_c = L_j(t_{\text{bo}} - \frac{r_{\text{bo}}}{c}) \sim L_jt_{\text{bo}} \] (13)
is the cocoon energy. Combining the above, we obtain an expression for the cocoon pressure,
\[ p_c = \frac{L_jv_{w} \delta t}{3r^3 \Omega_c (v_{jh} - v_w)}. \] (14)

The ram pressure of the wind material on the cocoon is given by \( p_{\text{ram,wc}} = p_w v_w^2 \), where \( v_w^2 = \Omega_c r^2/(\pi \delta t) \) is the velocity squared of the shock front driven by the cocoon into the ejecta. We obtain:
\[ p_{\text{ram,wc}} = \frac{m_w \Omega_c (v_{jh} - v_w)^2}{\pi \Omega_c v_w^2 \delta t^2}. \] (15)

Finally, let us consider the jet pressure, which has two components. First, the internal pressure of the jet, which is a relativistic invariant, and given by:
\[ p_j = \frac{L_j^2}{4\Omega_j^2 r^4 v_{j,inj}^2}. \] (16)

Additionally, one must consider the ram pressure of the jet material that is deflected from its initially radial velocity into a cylindrical flow in which the velocity is predominantly axial.

This pressure component is obtained by correcting the jet ram pressure with a geometrical factor \( \sin^2 \theta_{j,\text{inj}} \), yielding
\[ p_{\text{ram,jc}} = \frac{L_j \Omega_j^2 \sin^2 \theta_{j,\text{inj}}}{\Omega_c r^2}. \] (17)

In most cases, unless the jet is not significantly recollimated, the ram pressure dominates and hence we will use Equation (17) for the jet pressure on the cocoon hereafter.

Using Equations (14) and (15) in Equation (9) we obtain an equation for the cocoon solid angle as
\[ \Omega_c^2 = \frac{\pi L_j v_w^4 \delta t^3 \Omega_w}{3r^3 (v_{jh} - v_w)^2 m_w}. \] (18)
Alternatively, an expression for the cocoon solid angle can be derived using Equations (14) and (17) in Equation (8) to yield
\[ \Omega_c = \frac{\Omega_c v_w \delta t}{3r (v_{jh} - v_w) \sin^2 \theta_{j,\text{inj}}}. \] (19)

Equating Equation (18) to the square of (19) and solving for the jet solid angle we obtain:
\[ \Omega_j = \frac{3\pi L_j v_w \Omega_w}{m_w c^2 v_{jh}} \sin^2 \theta_{j,\text{inj}} \sin^2 \theta_{j,\text{inj}} \] (20)

The right-hand side of this equation can now be inserted for the jet opening angle in Equation (12) and solved for the jet head velocity, by using Equation (6) for the radius \( r \).

Finally, the cocoon energy can be derived from Equation (13) and, thanks to energy conservation, the remaining jet energy \( E_j = E_{j,\text{inj}} - E_c \), where \( E_{j,\text{inj}} \) is the energy injected at the base of the jet.

To validate this model we show in Figure 1 a comparison between the numerical results of Lazzati et al. (2017) and the energy and gamma profiles derived from the above equations.
A double exponential fit (orange dashed line) to the numerical results (blue solid line) yields $\theta_0 = 12^\circ 8$ and $\theta_c = 12^\circ 79$, respectively. Our analytic model (green solid line) yields instead $\theta_0 = 32^\circ 5$ and $\theta_c = 18^\circ 7$. While there are some differences due to the complexity of the interaction between a relativistic jet and the relatively slow and dense ejecta, the main features are well represented in the analytic model. We stress that the simulation was used as an input for the functional form in Equations (1) and (2), but the scale angles and cocoon energy used to draw the green curve in the figure are obtained from the analytic model. The general agreement between the numerical result (in blue) and the model (in green) is therefore not a circular argument.

Figure 2 shows the jet and cocoon opening angles (top and middle panels), and cocoon energies (bottom panels) as a function of the total mass in the ejecta at the time of the jet injection and of the delay between the launching of the ejecta and the release of the relativistic jets. From left to right, ejecta velocities $v_w = 0.1 c$, $0.3 c$, and $0.5 c$ are shown.

Figure 3 shows the energy (top panels) and Lorentz factor (bottom panels) profiles for a few selected cases. The cases shown in Figure 3 are indicated in Figure 2 with colored star symbols.

3. Afterglow Light Curves

We compute the afterglow light curves by feeding the energy and Lorentz factor profiles into a semi-analytic external shock code. We adopt the same version of the radiative code that was previously used by Lazzati et al. (2017a, 2017b, 2018) and Perna et al. (2019). The code uses an analytic approximation for the $n(t)$ profile that correctly approaches asymptotically the relativistic limit ($r \propto t^{1/4}$) at early times and the Sedov–Taylor scaling ($r \propto t^{1/3}$) at late times. However, throughout the evolution the code assumes a Blandford–McKee self-similar profile (Blandford & McKee 1976) for the blast wave structure behind the shock. This yields an approximate evolution for the afterglow, especially at the later stages, when the transition to an NR expansion alters the downstream shock profile (e.g., Huang et al. 1999, 2000; De Colle et al. 2012; Pe’er 2012; Li et al. 2019). The light curves shown here, when compared with more accurate calculations such as those used in BoxFit (van Eerten et al. 2012), are correct within $\sim$20% even at late times.

The external medium is assumed to be a low-density, uniform interstellar medium of density $n_{\text{ISM}}$, and the shock microphysics is parameterized, as customary, with the equipartition parameters $\varepsilon_e$ and $\varepsilon_B$, and with the nonthermal electron’s energy distribution index $p_{\text{el}}$. All electrons are assumed to be accelerated into the nonthermal distribution. We perform our light-curve calculations adopting typical average values for the shock microphysics parameters, as inferred in the afterglow modeling of the Swift-detected short gamma-ray bursts (GRBs; Fong et al. 2015) and of GRB170817 (Lazzati et al. 2017a). These are $\varepsilon_e = 0.1$, $\varepsilon_B = 0.01$, and $p_{\text{el}} = 2.3$. The interstellar medium is assumed to be uniform and tenuous, as expected in the surrounding of a BNS merger, with density $n_{\text{ISM}} = 10^{-2}$ cm$^{-3}$. The observer angle with respect to the jet axis is set to $\theta_{\text{jet}} = 30^\circ$, the typical viewing angle expected for LIGO/Virgo detected events (Schutz 2011) and the distance to the burst is set to 40 Mpc as a representative case to relate to GRB170817.

Examples of light curves in X-ray (1 keV), optical (R-band), and radio (6 GHz) are shown in the top and bottom panels of Figure 4, respectively. Interestingly, the light curves are noticeably different from one another, giving hope that the ejecta mass and velocity, as well as the time delay $\delta t$, can be constrained with a high quality data set. In particular, the slope of the slow-rising phase before the maximum and the prominence of the feature at the peak are characteristic of the jet–cocoon configuration. A larger ejecta mass results in a more energetic cocoon confined in a narrower scale angle $\theta_c$ (Figure 2) that moves with lower Lorentz factor (Figure 3). The lower Lorentz factors cause a late onset of the afterglow at both X-ray and radio frequencies, which is therefore a signature of significant mass ejection. Large ejecta masses also cause the jet scale angle $\theta_0$ to narrow significantly, resulting in a clearer jet feature at the light-curve peak, especially at radio frequencies (see, e.g., the green light curve in the top central panel of Figure 4). Very-low-mass ejecta produce instead a very weak but fast-moving cocoon and wide angle jet. These light curves peak early, especially in X-rays, and more closely resemble those of canonical, on-axis jets. Note that in the complete absence of ejecta, the light curves peak late, since no cocoon is present, as shown with the case of a pure jet by Perna et al. (2019). This might seem contradictory, but there is indeed a big difference between the complete absence of ejecta, which leaves the top-hat jet unaltered, and the presence of light ejecta, which produce a fast, low-energy cocoon. The former configuration gives a peak time that corresponds with the time at which the jet emission comes into the line of sight, at about 100 s in Figure 4. The latter, instead, peaks early, since the deceleration radius scales as $r_{\text{dec}} \propto E^{1/3} / \Gamma_c^{-2/3}$, yielding a deceleration time for material along the line of sight $t_{\text{dec}} \propto E^{1/3} / \Gamma_c^{-8/3}$. This yields a scaling for the cocoon-afterglow peak luminosity $L_{\text{pk}} \propto E^{2/3} / \Gamma_c^{8/3}$, showing that the peak luminosity of the cocoon grows quickly for increasing $\Gamma_c$, unless its energy vanishes.

In some cases, however, there are some degeneracies. For example, a careful comparison of the leftmost and rightmost panels of Figure 4 reveals similar light curves at both radio and X-ray frequencies. This is not surprising since both the delay $\delta t$ and the ejecta velocity $v_w$ control the radius of the ejecta in which the jet is propagating. Breaking these degeneracies will require the use of additional data. For example, one could use the shock-breakout model of Kasliwal et al. (2017) to constrain the outer radius of the ejecta from the brightness of the prompt $\gamma$-ray emission. In addition, constraints on $\delta t$ can be set since it has to be smaller than the time delay between the detection of GWs and of $\gamma$-ray photons. Additional constraints could come from the kilonova properties and from the GW detection (e.g., constraints on the inclination of the system, the chirp mass, and the individual masses of the two NSs). Thus one can envisage that afterglow modeling with the appropriate priors from the complete data set will be able to set meaningful constraints on the ejecta properties.

4. Ejecta Properties and the Equation of State of NSs

As already mentioned in Section 1 to motivate this work, the connection between the amount and velocity of the ejected mass and the observable properties of the mergers (in particular, in the electromagnetic spectrum) is especially interesting in light of the fact that it contains information on the EoS of the NS, as demonstrated by several groups via GRMHD simulations (e.g., Shibata et al. 2005; Kiuchi et al. 2010; Rezzolla et al. 2010; Hotokozzaka et al. 2011, 2013;
Most of the tidally disrupted mass remains bound and circularizes into a hot torus, whose accretion onto the remnant compact object formed from the merger powers a relativistic jet. The total energy in the jet, proportional to the total accreted mass, carries information on the EoS; this connection was exploited by Giacomazzo et al. (2013) with the sample of the short GRBs available at the time, to draw inferences on the likelihood of various EoS.

Here we are focusing on signatures of the ejecta mass in the light curves. This high-velocity material is made of several components. First, the dynamical ejecta (e.g., Bauswein et al. 2013; Giacomazzo & Perna 2013; Kawamura et al. 2016; Lehner et al. 2016; Ruiz et al. 2016; Sekiguchi et al. 2016; Radice et al. 2016, 2018b; Ciolfi et al. 2017; Dietrich et al. 2017; Foucart et al. 2019).\(^4\)

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which is partly due to matter ejected (and unbound) by the tidal disruption of the less massive NS, and partly by matter ejected (and unbound) during merger due to the formation of shocks. More material is expected to be ejected via ablation from the surface of a hypermassive NS (HMNS) in the case in which the merger remnant passes through such a phase (e.g., Dessart et al. 2009; Perego et al. 2014, 2017; Just et al. 2015). Additionally, nuclear and viscous processes in the accretion disk surrounding the remnant compact object unbound material from the disk, creating a disk wind (e.g., Metzger et al. 2009; Fujibayashi et al. 2017, 2018; Fahlman & Fernández 2018; Radice et al. 2018b; Siegel & Metzger 2018), also referred to as secular ejecta. Recent simulations (e.g., Siegel & Metzger 2017; Radice et al. 2018b) find that the secular ejecta mass typically exceeds the mass of the dynamical ejecta, except for the cases of prompt BH formation, in which they can be comparable (albeit small). The quantitative results for the masses of both components are still rather uncertain, as they depend on a number of microphysical effects not all of which are fully included at this stage in the numerical simulations. Neutrino cooling, for example, plays an important role, and so does the magnitude of the effective viscosity arising from magneto-hydrodynamic instabilities operating during the merger. For example, Radice et al. (2018a) found that the mass of the fast tail of the dynamical ejecta can vary by up to four orders of magnitude depending on the viscosity strength.

In terms of relevance for our problem the ejecta mass and velocity, their launch time with respect to the jet, and their angular distribution play a role. Fernández & Metzger (2013) showed that disk outflows can develop only after the disk becomes advective and $\alpha$-particles form. They find that the mass flux takes about one second before reaching its maximum. If the disk wind is in fact delayed with respect to the jet launch, then the jet will mostly interact with the dynamical ejecta. What matters for this work is the fact that both the dynamical ejecta and the wind outflow carry the imprint of the EoS of the NS, in addition to being dependent on the masses and mass ratios of the NSs. Generally speaking, the simulations show that larger mass ratios tend to result in the partial disruption of the smaller star during the merger, and hence a smaller amount of shocked material, since the stars merge less violently. The mass asymmetry also produces more massive tidal outflows. A larger amount of ejecta further tends to correlate with softer EoSs.

A visual representation of the dependence of the ejecta properties on the EOS of NS is offered in Figure 5, which displays the dynamical ejecta mass (top panels) and velocity (bottom panels) for three representative equations of state of NS matter. These quantities are shown as a function of the

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*Figure 3. Energy profiles (top panels) and Lorentz factor profiles (bottom panels) for the cases indicated with colored star symbols in Figure 2 (see also the figure legend).*

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Data from https://compose.obspm.fr/.
mass ratio $q$ and the chirp mass, which is extremely well constrained by GW observations alone (e.g., The LIGO Scientific Collaboration et al. 2018). In the figures, white areas correspond to regions that are either unphysical or for which the mass–radius relationship is not allowed by that equation of state. The plots of Figure 5 were made using the fits by Coughlin et al. (2018a), based on an update of the ones by Dietrich & Ujevic (2017) with the inclusion of more simulations (for a total of 259); however, they describe only the dynamical component of the ejecta, for which a much larger number of simulations (and parameter exploration) exists in the literature to allow the derivation of fitting formulae. For the disk wind and the HMNS wind component, fewer simulations exist. Some fits to the ejecta mass have been provided in the literature (e.g., Coughlin et al. 2018b; Radice et al. 2018b), but no systematic fit exists to date for the ejecta speed (see, however, some correlations with the disk mass provided by Fahlman & Fernández 2018). Therefore, the connection between ejecta properties and EoS will need to be further tightened in the future with more accurate simulations. At the same time, it will also be important to better understand the time of jet launch with respect to the other outflow timescales in the problem.

5. Summary and Conclusions

The detection of electromagnetic radiation from a binary NS merger triggered by GWs has allowed an unprecedented view into the inner workings of a relativistic jet, likely powered by accretion onto the remnant object, and propagating onto the ejecta from the merger.

Since the LIGO horizon is much smaller than the Swift one, events triggered in GWs will be generally much closer than the ones triggered in $\gamma$-rays. Hence their afterglows, especially in the radio band, which is not contaminated by the kilonova emission, will generally be much brighter than for the more distant short GRBs triggered in $\gamma$-rays. As shown by the case of GW170817, this allows for a detailed monitoring of the afterglow light curves in several bands by a variety of instruments, and hence a detailed shape reconstruction of the light curve.

In this paper we have studied the structure of the jet–cocoon, which forms as a result of the jet expanding into the fast-moving ejecta produced by the NS–NS merger. More specifically, we have explored how this structure varies with the density, velocity, and basal extent of the ejecta. We have then used this structure at the radiative stage to compute the shape of the afterglow in the X-ray and radio bands (the optical
is likely dominated by the kilonova, hence may have less diagnostic power, and explored its dependence on the ejecta properties, as well as on the time delay between merger and onset of the jet.

We identified features in the light curves, which are especially distinctive of a jet–cocoon configuration, where the cocoon is generated by the interaction of the jet with the dynamical ejecta. A more pronounced cocoon yields an additional, earlier bump or plateau in the light curve, compared to the singly peaked light curve of a pure relativistic jet. The rise time of the light curves is also very sensitive to the amount of ejecta, with larger values yielding delayed rises. It should be noted that our scenario might be oversimplified in some aspects and further observations or advancement in the understanding of the ejecta of BNS mergers could spark the need for modifications. In particular, the ejecta might not be spherically (or even cylindrically) symmetric. An oblate configuration, for example, would cause more efficient collimation of the jet (e.g., Duffell et al. 2015). In addition, the central engine might not become dormant after a few seconds, as assumed here. A long-active central engine, such as a magnetar, may result in extra energy injection in both the jet (Geng et al. 2018) and the ejecta (Yu et al. 2018), altering the simple geometry and dynamics that we assumed.

In conclusion, a detailed monitoring of the afterglow light curve, especially in the cleaner X-ray and radio bands, can help constrain the properties of the ejecta from the merger event. Since this material is sensitive to the equation of state of dense matter (in addition to being important for production of heavy elements), detailed monitoring of afterglow light curves of GW-detected binary NS mergers (and NS–BH alike) has the power to help probe the physics of the merger events and the equation of state of dense nuclear matter.

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