A cocoon shock breakout as the origin of the $\gamma$-ray emission in GW170817

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ABSTRACT

The short Gamma-Ray Burst, GRB170817A, that followed the binary neutron star merger gravitational waves signal, GW170817, is not a usual sGRB. It is weaker by three orders of magnitude than the weakest sGRB seen before and its spectra, showing a hard early signal followed by a softer thermal spectrum, is unique. We show, first, that the $\gamma$-rays must have emerged from at least mildly relativistic outflow, implying that a relativistic jet was launched following the merger. We then show that the observations are consistent with the predictions of a mildly relativistic shock breakout: a minute $\gamma$-ray energy as compared with the total energy and a rather smooth light curve with a hard to soft evolution. We present here a novel analytic study and detailed numerical 2D and 3D relativistic hydrodynamic and radiation simulations that support the picture in which the observed $\gamma$-rays arose from a shock breakout of a cocoon from the merger’s ejecta (Kasliwal et al. 2017). The cocoon can be formed by either a chocked jet which does not generate a sGRB (in any direction) or by a successful jet which generates an undetected regular sGRB along the system’s axis pointing away from us. Remarkably, for the chocked jet model, the macronova signal produced by the ejecta (which is partially boosted to high velocities by the cocoon’s shock) and the radio that is produced by the interaction of the shocked cocoon material with the surrounding matter, agree with the observed UV/optical/IR emission and with current radio observations. Finally, we discuss the possibility that the jet propagation within the ejecta may photodissociate some of of the heavy elements and may affection the composition of a fraction of ejecta and, in turn, the opacity and the early macronova light.

Key words: gamma-ray burst: short | stars: neutron | gravitational waves | methods: numerical

1 INTRODUCTION

The advanced LIGO and advanced Virgo gravitational radiation observatories have detected the first binary neutron star merger on August 17 2017 (The LIGO Scientific Collaboration, The Virgo Collaboration 2017a,b,c,d). The gravitational radiation event known as GW170817 revealed a merger of two low mass compact objects, whose mass range clearly puts them as neutron star rather than black hole candidates (The LIGO Scientific Collaboration, The Virgo Collaboration 2017d; B. P. Abbott 2017). Remarkably, the $\gamma$-ray satellite Fermi detected a short gamma-burst (GRB) about 2 seconds after the GW signal, from a location that is consistent with the localization of the GW signal (Goldstein 2017). This was followed by a detection of an optical counterpart, SSS17a (Coulter 2017), associated with the accompanying Macronova/kilonova that was detected also in the UV in the IR (Evans et al. 2017) and later on with an X-ray (Troja et al. 2017) and Radio (Hallinan et al. 2017) counterparts. The optical observations allowed the identification of a host galaxy, NGC 4993, an early-type galaxy at a distance of $\approx 40$ Mpc.

While tantalizing and seemingly confirming a long standing prediction of the association of sGRBs with mergers (Eichler et al. 1989), a quick look at the $\gamma$-ray observations shows that the sGRB 170817A which was discovered in association with GW170817 is not a regular sGRB. While its fluence, $(2.8 \pm 0.2) \times 10^{-7}$ erg cm$^{-2}$, and duration, $2 \pm 0.5$ s, are similar to other observed sGRBs, when taking into account its distance of $\sim 40$ Mpc, we find that its total isotropic equivalent energy, $E_{iso} = (5.35 \pm 1.26) \times 10^{46}$ erg is smaller by three orders of magnitude than the weakest sGRB whose energy was measured so far and by four orders of magnitude than typical sGRBs (see e.g. Nakar 2007, for a review). Furthermore, the pulse is made of two parts with a very different spectrum. The spectrum of the main pulse ($T_0-0.320$ s to $T_0+0.256$ s) is best fit with a Comptonized function, with a power-law index of $-0.62 \pm 0.40$ with peak energy $E_{peak} = 185 \pm 62$ keV. The main pulse is followed by a weaker tail extending from $T_0+0.832$ s to $T_0+1.984$
was discovered here is a new type of gamma-ray bursts - a low luminosity short GRB, denoted hereafter ltsGRB.

Compactness arguments (Kasliwal et al. 2017) reveal that the observed γ-rays must have been produced in a mildly or fully relativistic outflow with Γ ≥ 2.5. At the same time Kasliwal et al. (2017) showed that three simple scenarios that come to one mind: a very weak relativistic jet pointing towards us; an off-axis emission of a strong GRB pointing along the rotation axis (30° away from us); a structured jet with a broad weak wing, are highly unlikely. We must therefore turn to something else. As we show in §2 the fact that the emitting region must be moving relativistically cannot arise in a spherical explosion and therefore the merger must involve a relativistic jet. Numerical simulations have shown long ago (e.g. Davies et al. 1994; Ruffert et al. 1996; Rosswog et al. 1999; Bauswein et al. 2013; Hotokezaka et al. 2013; Sekiguchi et al. 2015; Radice et al. 2016) that neutron star mergers are accompanied by a significant amount of mass ejection. The observations of the monacrona accompanying GW170817 confirms this idea putting a lower limit of 0.02M⊙ on the mass ejected in this event (Kasliwal et al. 2017). Any relativistic jet will have to path through this mass (see Nagakura et al. 2014; Murguia-Berthier et al. 2014; Duffell et al. 2015; Lazzati et al. 2017a; Gottlieb et al. 2017 for jet propagation studies). Regardless of the question whether the jet penetrates the ejecta or is choked within it, it will produce an energetic cocoon. We have shown in the past that the cocoon can generate a mildly relativistic outflow at a wide angle (Gottlieb et al. 2017). Here we suggest that the observed γ-rays are produced when the shock driven by this mildly relativistic cocoon breaks out of the ejecta. This process is drastically different from the one that takes place in regular sGRBs. In fact, we suggest that there are similarities between the physical mechanisms (but not in the astronomical scenario) deriving this ltsGRB and those that take place in regular (long) low-luminosity GRBs (llGRBs).

A shock breakout of the cocoon provides a natural explanation to two puzzling properties of the observed pulse. First, the energy emitted in γ-rays in our direction is only a minute fraction, ~10⁻⁴ (isotropic equivalent) of the total kinetic energy seen in the outflow (~10⁵¹ erg). Although a jet that can propagate a significant way through the ejecta, and thus any cocoon that it produces, is expected to carry much more energy than the one we observe in γ-rays. This is unlike regular GRBs, which are very efficient in γ-ray production. However, it is a natural property of shock breakout emission, where the radiation is coming from a very narrow layer where the optical depth upon breakout is low enough, while the bulk of the energy is in highly optically thick material and is therefore hidden (this energy is released only much later after it suffers significant adiabatic losses due to expansion). Similar bulk to breakout energy ratios are seen in breakout candidates such as SN2008D (Soderberg et al. 2008) and low-luminosity long GRBs (e.g., Kulkarni et al. 1998, ?). Second, one of the predictions of a relativistic shock breakout from a star is an initial hard pulse in γ-rays, followed by a softer (typically X-ray) tail that carries a comparable amount of energy (Nakar & Sari 2012). Finally, the model we present here reproduces all the main features of the observed γ-rays, the total energy, the duration, the spectrum (and its variation with time) and the 2 s time delay between the gravitational waves and the γ-rays.

We present here a new order-of-magnitude analytic model for a shock breakout from an expanding ejecta and highlight the differences from a shock breakout from a stellar envelope. We then carry out the first detailed numerical calculations of a cocoon shock breakout in sGRBs, including the time evolution of both the luminosity and the spectrum, as well as the dependence on the viewing angle. The structure of the paper is as follows. We begin in §2 demonstrating using compactness that, although we did not observe a regular sGRB, the emitting region must have had a relativistic jet. We discuss our model in §3. Our numerical setup, which deals with the hydrodynamic of an unmagnetized jet, and the calculation method of the breakout signal are presented in §4. A companion paper (Bromberg et al. 2017) discusses relativistic MHD simulations for the propagation of a magnetic dominated jet. We present the configuration of a choked jet and the resulting γ-rays, UV/optical/IR and radio signals in §5. A shock breakout from a cocoon produced by a successful jet, that emerges and powers a regular sGRB pointing away from us, can also produce the observed γ-rays. We describe such a system in §6. Before concluding we briefly discuss in §7 the possibility of dissociation of nuclei within the shocked ejecta matter via photodissociation or due to interaction with relativistic neutrons. We show that in the setups we have considered these effects can most likely be ignored, but that it may be important in other scenarios or other events. Detailed calculations of this process will be discussed elsewhere. We conclude with implications and observational predictions in §8.

2 COMPACTNESS AND RELATIVISTIC MOTION

Remarkably, in spite of the weakness of this burst its γ-ray emission clearly demonstrates that the first non-thermal component of this burst must have involved relativistic motion towards us. The simple compactness arguments show that if the source would have been Newtonian, its optical depth due to pair production would have been too large to be consistent with the observations. The optical depth of the emitting region can be estimated as (Nakar 2007; Kasliwal et al. 2017):

\[ \tau \approx \frac{3\sigma_{\gamma}\Gamma}{E^{\gamma}A(c\tau_{0.1})^2} \approx 6 \times 10^{57} f(\Gamma) \frac{E_{100}}{3 \times 10^{46} \text{ erg}} \left( \frac{E_{\gamma}}{150 \text{ keV}} \right)^{-1} \left( \frac{T_{0.1}}{0.6 \text{ s}} \right)^{-2} \]

where Γ is the Lorentz factor of the emitting region E_{100} is the isotropic equivalent energy of the first component, E_{\gamma} is the average energy of its photons, T_{0.1} its duration and f(\Gamma) is the fraction of photons of this component that are above the e^+e^- production threshold. In GRB 170817A this is a dominant factor as the spectrum drops exponentially f(\Gamma) ~ \Gamma^{-0.6} \exp[-(\Gamma m_{e}c^2)/125\text{keV}]. Now, if the source is Newtonian then τ ~ 10^6, implying first a diffusion time that is too long unless the emission comes from an unrealistically thin shell (<10^5 cm while the emitting region is at radius ~10^{10} cm) and even if such a shell existed the extremely high density of pairs in it (~10^{26} cm^-3) would have generated enough photons during the diffusion time to obtain thermal equilibrium at a temperature much lower than the observed E_{\gamma}.

Thus the source cannot be Newtonian. However the exponential cut-off in the spectrum makes f(\Gamma) extremely sensitive to the Lorentz factor, and therefore even a mildly relativistic Lorentz factor of 2.5 is enough to make the source optically thin. We therefore conclude that the emitting region must have been moving with at least a mildly relativistic velocity. An immediate implication of this result is that a relativistic jet must have been involved. It is clear
from the optical observations that the bulk of the outflow is moving at 0.1-0.2c (Kasliwal et al. 2017). This is much slower than the required speed. The spherical ejecta can have a fast tail, and in fact we argue later that there is one (with \( \sim 10^{-7} M_c \) moving at a typical velocity of \( \sim 0.6-0.8c \)). However, the highest velocity component of the spherical ejecta material is ejected over a very short time scale \( \sim 10 \) ms and at a radius \( \lesssim 10^7 \) cm, while the duration and the delay of the \( \gamma \)-ray signal implies that it should have been radiated at a radius of \( \sim 10^{11} \) cm. At this radius the ejecta has already lost all its internal energy due to adiabatic expansion and it moves homologously. Therefore, while this material contains enough energy to power the observed \( \gamma \)-rays, there is no internal mechanism to this outflow which can dissipate the bulk kinetic energy to internal energy and produce \( \gamma \)-rays at this point. The most plausible way to dissipate this energy is via energy emitted from the post merger compact object. There is \( \gtrsim 0.02 M_c \) between the central engine and the ejecta front, which are moving at \( \lesssim 0.2c \) also after the energy that generates the \( \gamma \)-rays is deposited in the outer layers. Therefore, this energy cannot be channeled out by a spherical outflow and a relativistic jet seems to be necessary.

### 3 A RELATIVISTIC SHOCK BREAKOUT FROM AN EXPANDING EJECTA

As mentioned earlier, the optical/IR observations (Kasliwal et al. 2017) clearly indicate the existence of a few percent of a solar mass that have been ejected during this merger. Most of this material is moving at 0.1-0.2c. A regular sGRB jet would have had to penetrate through this ejecta before emerging and producing the \( \gamma \)-rays. A second possibility is that the central engine did produce a jet but this jet was choked as it could not penetrate through the ejecta (either due to a low power, short duration or a large opening angle). In fact, just a short while ago Moharana & Piran (2017) inspected the duration distribution of sGRBs and suggested an evidence for a significant number of sGRBs with choked jets. Gottlieb et al. (2017) have shown that a jet that successfully penetrates the ejecta can drive a mildly relativistic cocoon which expands over a large angle. Here we show that a choked jet is capable of doing that as well, if it is powerful enough. Thus, both scenarios generate a mildly relativistic shock breakout from the expanding ejecta. Below we discuss the properties of the signal that such a breakout produces.

The jet drives a radiation mediated shock into the expanding ejecta. The jet drives a radiation mediated shock into the expanding ejecta. Since the jet is relativistic and it is expanding into a rather dilute plasma its head velocity is at least mildly-relativistic and hence the shock is also mildly relativistic, regardless of the details of the density profile.

The shock propagates as long as the optical depth to infinity, \( \tau \), is large enough to sustain its width. Once \( \tau \) drops below this point, the radiation in the shock transition layer escapes to infinity and the shock ‘breaks out’. Following the breakout the radiation that was deposited by the shock starts diffusing out of the expanding gas, generating the so called ‘cooling emission’. The theory of shock breakout and the emission that follows was studied mostly in the context of a shock that propagates in a stellar envelope, both at Newtonian velocities (e.g. Colgate 1974; Chevalier 1976; Falk 1978; Matzner & McKee 1999; Nakar & Sari 2010) and relativistic velocities (Nakar & Sari 2012).

Breakout from an expanding ejecta with an edge, as expected in this case, was not considered so far in the literature. It is different than a breakout from a static stellar envelope in two important ways. First, there are two characteristic velocities instead of one, the shock velocity as seen in the lab frame, which determines the boost of the emission to the observer, and the shock velocity in the upstream (i.e., ejecta) frame, which determines the shock strength and can be significantly different from the former. Second, unlike the breakout from a stellar edge, the diffusion length scale (i.e., the scale over which the optical depth changes significantly) and the hydrodynamic length scale (i.e., the scale over which the density, pressure and velocity change significantly) can be different. As explained below these differences can have an important effect on the observed emission.

The resulting light curve and spectrum can depend on details such as the ejecta profile and the density near its edge. We defer a detailed study of the signal to a future study. Here we discuss first the dynamical evolution which is common to all breakout scenarios followed by a discussion of the spectrum that this evolution dictates. We then discuss the main robust features that the general dynamical evolution induce for a wide range of the breakout parameters. We also derive an order of magnitude estimate of the breakout parameters that are needed in order to produce the observed signal, showing that such parameters exist, that they are over constrained (i.e., more observables than free parameters) and very reasonable.

#### 3.1 Planar and spherical phases

We consider a breakout of a spherical relativistic shock, with a Lorentz factor \( \Gamma \), and velocity \( v = \beta c \) as seen in the lab frame, from ejecta that expands with a maximal velocity \( \beta_{e,j,max} \). that can be Newtonian or mildly relativistic. The shock velocity as seen in the ejecta frame, \( \beta' = (\beta - \beta_{e,j,max})/(1 - \beta \beta_{e,j,max}) \), can be mildly relativistic or even sub-relativistic. The shock breakout takes place at a radius \( R_{bo} \) from a layer with an optical depth \( \sim c/\sqrt{\kappa_{e,j}} \), which we denote as the "breakout layer." In general this layer is much narrower in the lab frame than the causality scale \( R_{bo}/\Gamma^2 \). Therefore, the hydrodynamics, and as a result the observed signal, has two phases: planar and spherical. The planar phase starts right after the breakout and it lasts a duration \( \sim R_{bo}/c \) in the lab frame (\( \sim R_{bo}/c \Gamma^2 \) in the observer frame), namely until the breakout layer doubles its radius. During this phase the radiation in the breakout layer, as well as in layers with higher optical depth which may be important (as we show later), evolve on time scales much shorter than \( R_{bo}/c \). During this phase the radius can be approximated as a constant. In the spherical phase everything evolves on a single time scale, \( R_{bo}/c \) in the lab frame (however, see also Yalinewich & Sari 2017 for the case of an acceleration in the spherical phase). Below we discuss shortly each of these phases.

**Planar phase:** The difference between a breakout from a star and a breakout from a relativistic ejecta is most important in the planar phase. In a star, during this phase, a fluid element has only a single length scale, the initial distance from the stellar edge. Consequently, the hydrodynamical scale and the diffusion scale are similar and the time it takes the breakout shell to expand is similar to the time it takes the radiation to diffuse through the breakout layer. This implies that during the entire planar phase only the breakout shell releases radiation to the observer (see Nakar & Sari 2010 for a detailed discussion). In an expanding ejecta, the width of the breakout layer can be much smaller than the hydrodynamical scale which we expect to be of order \( R_{bo}/\Gamma \) in the fluid rest frame. In this case the breakout layer does not expand significantly over a diffusion time scale and photons from inner layers have time to diffuse out during the planar phase. This has two effects: First, it increases the...
luminosity of the planar phase and second it modifies the spectrum since, as discussed below, the temperature of photons drops with the time it takes them to diffuse, and thus the spectrum of the planar phase is non-thermal, composed of the sum of radiation of different temperatures. Note that due to light-travel-time effects the emission from the planar phases of all the contributing layers is smeared over the angular time and seen together over a duration of \( \sim R_{bo}/2c\Gamma_{s}^2 \).

During the spherical phase the radius increases, the optical depth decreases and emission from inner layers diffuses out more efficiently. Photons from the spherical phase are emitted after the expanding gas doubled its radius, namely they are observed at a time \( \sim R_{bo}/2c\Gamma_{s}^2 \) after the breakout emission. Therefore, they are not mixed with the planar phase photons. The layers that radiate during the spherical phase have much more time to cool and get closer to thermal equilibrium (see below). They also go through some adiabatic expansion and cooling. In addition, the angular time and the dynamical time are now comparable, implying that angular smearing does not mix much radiation with different temperatures. Therefore the spectrum at the beginning of the spherical phase is expected to be much softer than in the planar phase, and it is much more well defined by a single temperature. The spherical phase, that is often referred to as the cooling emission phase, continues for a long time after the breakout. During this phase both the luminosity and the temperature drop with time. The luminosity drops gradually, while the temperature may drop quickly at first as the emitting layers are out of thermal equilibrium at the beginning of the spherical phase (see Nakar & Sari 2010 for details).

The planar and spherical phases are well separated in the observed signal when the breakout is spherical, as assumed above. If the breakout is oblique and takes place at different radii for different angles, then angular smearing can mix the emission of planar phase at one angle with the emission of spherical phase from another. If this smearing is strong then the spectrum still shows a hard to soft evolution, possibly with two well separated components, but it will remain non-thermal during the entire evolution.

### 3.2 The observed spectrum

The observed spectrum depends on the radiation temperature in the diffusing gas. When the shock is fast enough (\( \gtrsim 0.05c \) in the upstream frame) the radiation behind the shock falls out of thermal equilibrium (Weaver 1976; Katz et al. 2010; Nakar & Sari 2010). The temperature \( T \) is determined then by the ability of the gas to generate photons during the available time, which becomes less efficient when the shock is faster. As a result the radiation temperature behind the shock rises sharply with the shock velocity, and at \( v_{s} \approx 0.5c \) it reaches \( \sim 50 \) keV at which pair production becomes important. Once pairs are produced efficiently the photon production rate rises as well as the exponential dependence of the number of pairs on the temperature serves as a thermostat that sets the temperature behind the shock at \( \sim 100 \sim 200 \) keV for a relativistic shock, regardless of its Lorentz factor (Katz et al. 2010; Budnik et al. 2010; Nakar & Sari 2012; Granot et al. 2017).

The temperature evolution after the breakout depends on the pair loading of the shock. If the shock is not too fast (\( v_{s} \lesssim 0.5c \)) and pairs are negligible then the radiation is released upon the breakout and the observed temperature is the gas temperature behind the shock times its Lorentz factor as seen in the observer frame, \( \sim \Gamma_{s} \).

If the shock is loaded by pairs (\( v_{s} \gtrsim 0.5c \)) then the radiation remains trapped until its rest frame temperature drops to \( \sim 50 \) keV. At this stage the pairs annihilate and the photons are released. In that case the observed temperature is \( \sim 50 \) keV times the Lorentz factor of the radiating gas (Nakar & Sari 2012).

The radiation that diffuses to the observer following the shock breakout spends more time trapped in the gas and therefore the gas has more time to generate photons to share the internal energy, thereby reducing the radiation temperature. Quantitatively the temperature of the diffusing radiation can be estimated using the approximation presented in Nakar & Sari (2010), that followed the derivation of Weaver (1976). This approximation solves for the gas temperature, \( T \), at which the number density of generated photons that can be coupled to the gas energy, \( n_{ph} \), is enough to share the entire gas energy density, \( \epsilon \), namely \( \epsilon = \eta_{ph}3k_{b}T \), where \( k_{b} \) is the Boltzmann constant. For that we first estimate the production rate of photons that can share the gas energy, namely photons that are emitted at the gas temperature \( T \) or that can be Comptonized to that temperature in the available time. Since the typical temperatures are \( \gtrsim 10 \) keV we assume that the gas is fully ionized. We consider photon production by free-free and bound free. We calculate the free-free photon production below and approximate the bound-free photon production rate to be comparable. The free-free production rate of photons with an energy \( \sim k_{b}T \) by a gas with a mixed composition of heavy nuclei is then (e.g., Nakar & Sari 2010; Sapir et al. 2013)

\[
\eta_{ph, ff} \approx 3.5 \times 10^{36} \rho^{2}T^{-0.5} \left( \frac{z}{A} \right)^{2} \text{cm}^{-3}\text{s}^{-1},
\]

where \( \rho \) and \( T \) are in c.g.s, \( z \) and \( A \) are the atomic and mass numbers. The brackets mark an average over the gas and therefore depend on its composition. Here we approximate \( \left( \frac{z}{A} \right)^{2} \approx \frac{A}{A_{max}} = 10 \) for r-process material. The minimal frequency, \( \nu_{min} \), from which photons that are emitted by free-free at \( \hbar\nu_{\max} \approx k_{b}T \) can be Comptonized to \( T \) on time, satisfies (Weaver 1976; Nakar & Sari 2010; Sapir et al. 2013)

\[
\nu_{\max} = \frac{k_{b}T}{\hbar\nu_{min}} \approx 500 \left( \frac{0.1}{10^{9} \text{gr cm}^{-3}} \right)^{-1/2} \left( \frac{T}{1 \text{keV}} \right)^{9/4} \left( \frac{A}{A_{max}} \right)^{1/2},
\]

where we approximate \( (A)/(z^2)^{1/2} = 5 \) for r-process material. This increases the number of photons that can share the gas energy by a factor (Nakar & Sari 2010)

\[
\xi \approx \frac{1}{2} [\ln[\nu_{\max}](1.6 + \ln[\nu_{\max}])].
\]

The temperature in the diffusing gas is then found by solving the implicit equation

\[
\epsilon = 6k_{b}T\eta_{ph, ff}3k_{b}T \xi t,
\]

where \( t \) is the time passed in the shocked gas rest frame since the crossing of the shock and the release of the photons to the observer. A factor of 2 on the r.h.s accounts for the bound-free photon production which we approximate to be comparable to free-free. This estimate ignores pairs, which is appropriate for \( T \lesssim 50 \) keV. If \( T > 50 \) keV then pairs prevent the diffusion of photons until the temperature drops to \( \sim 50 \) keV.

A comparison of the results obtained from equation 5 with the dynamical simulation of Sapir et al. (2013) shows that the temperature during the breakout from a stellar surface is accurate to within a factor of \( \sim 2 \sim 3 \).
3.3 General properties

The two phases described above define several general properties that are common to the breakout signal over a wide range of configurations, including both a breakout from an expanding ejecta and from a stellar envelope. First, the breakout layer and all other layers, if there are any, that radiate during the planar phase contain only a very small fraction of the total internal energy of the expanding gas. As a result the breakout pulse contains a very small fraction of the total explosion energy. Some examples are the emission from lGRBs in which the γ-ray energy is only $10^{-3} - 10^{-4}$ of the total explosion energy (e.g. Kulkarni et al. 1998; ?) and SN2008D in which the energy carried by hard X-rays is $\sim 10^{-5}$ of the total SN kinetic energy (Soderberg et al. 2008). Second, the light curve cannot be highly variable. It may have some structure, especially upon the transition from the planar to the spherical phase, but it cannot produce the high variability observed in most short and long GRBs (e.g., Nakar & Piran 2002b,a). The shape of the light curve upon the transition from the planar to the spherical phase depends on how different are the diffusion and the dynamical scales and on whether the breakout is fully spherical or if it takes place at slightly different radii at different angles. If the scales are similar and the breakout is spherical (such as in a spherical explosion in a star), the planar phase produces a well defined bright pulse that is followed by a fainter more slowly evolving spherical phase emission. Otherwise, the transition between these phases becomes smoother. Third, regardless of how smooth is the transition of the luminosity from the planar to the spherical phase, the spectrum changes significantly. The peak of the emission is characterized by a relatively hard spectrum that is a composition of emission at different temperatures, while upon the transition to the spherical phase the spectral peak energy drops significantly and the temperature becomes more well defined. For a more spherical breakout the spectral transition is sharper and the temperature after the transition is more well defined. An example of these properties can be seen in the detailed solution of a relativistic breakout from a static stellar envelope presented in Nakar & Sari (2012) and in the results of the numerical simulations that we present here.

3.4 Order of magnitude estimates

We are interested in finding an order of magnitude estimate of the properties of a shock breakout that can produce a signal that is consistent with the observed γ-rays of GRB 170817A. Namely, what are $\Gamma_b$, $\beta_{e,j,max}$, $R_{bo}$ and the mass which from photons diffuse during the planar phase, $m_{bo}$, that generate a $t_{bo,obs} \sim 0.5$ s long γ-ray pulse during the planar phase with an energy of $E_{bo} \sim 3 \times 10^{46}$ erg and a typical photon energy of $\sim 100 - 150$ keV at a delay of $\sim 2$ s with respect to the merger time. Upon the transition to the spherical phase the temperature should drop to $\sim 10$ keV. Note that altogether there are four breakout parameters and five observables, so the problem is over constrained and there is no guarantee that there is a viable solution.

The observed temperature implies that the breakout velocity in the observer frame cannot be neutron. Consequently the velocity in the ejecta rest frame is also at least $v_j' \gtrsim 0.5 c$. Thus, the observed temperature satisfies $T \sim 50\Gamma^2_{bo}$ keV implying that the shock must be mildly relativistic with

$$\Gamma_{bo} \approx 2 - 3.$$ (6)

Note that if the shock had been mildly relativistic also in the upstream flow (i.e., $v_j' \gtrsim 0.7 c$ ), the gas would have accelerated significantly after the crossing of the shock and before the photons are released, so $\Gamma_{bo}$ would have been significantly larger than $\Gamma_j$ (see Nakar & Sari 2012 for details). We assume that this is not the case and $\Gamma_j \approx \Gamma_{bo}$ and verify consistency later. This breakout Lorentz factor together with the duration determines the breakout radius

$$R_{bo} \approx 2 c t_{bo,obs} \Gamma^2_{bo} \approx 2 \times 10^{11} \text{ cm}.$$ (7)

This radius is about 10 light seconds and it takes the shock about 11 s to travel from the source to the breakout radius. Thus, in order to obtain the observed delay the jet should be launched about a second after the merger. The outermost part of the ejecta has 2 seconds (denoted by $\delta t$) longer to expand than $R_{bo}/c$. Thus, it must have a velocity

$$\beta_{e,j,max} \approx \frac{R_{bo}}{R_{bo} + c \delta t} \approx 0.7.$$ (8)

If the shock Lorentz factor is $\Gamma_j \approx 2 - 3$ then this ejecta velocity satisfies, consistently $v_j' \approx 0.5 c$. Pair production is marginal and the gas does not accelerate significantly after shock crossing, namely the assumption $\Gamma_j \approx \Gamma_{bo}$ is satisfied.

The parameters that we find above determine the energy released during the breakout with limited freedom which depends on the mass carried by the fast tail of the ejecta (i.e., at ejecta velocity $\gtrsim 0.5 - 0.6 c$), $m_{tail}$. The minimal tail mass that is needed for a breakout to take place at $R_{bo}$, with $\beta_j' \approx 0.5$ is

$$m_{bo,\min} \approx \frac{4 \pi R^2_{bo}}{k \rho_1'} \approx 4 \times 10^{-9} M_{\odot},$$ (9)

where we use $\kappa = 0.15 \text{ cm}^2/\text{g}$ as expected for fully ionized heavy elements. If $m_{tail} = m_{bo,\min}$ then the width of the breakout layer in its rest frame, after it is shocked, is $\sim R_{bo}/\Gamma_j$ and the dynamical time for its expansion is comparable to the photon diffusion time through it. Therefore only this layer radiates during the planar phase. If $m_{tail} > m_{bo,\min}$ then the mass which contributes to the emission during the planar phase, and thus to the energy in the initial γ-ray pulse, grows as $\sim m_{bo} = (m_{bo,\min} m_{tail})^{1/2}$. With a shock velocity of $\beta_j' \approx 0.5$ the internal energy in the shocked breakout mass is $\sim 0.2 m_{bo} c^2$, implying that the energy of the initial breakout pulse is:

$$E_{bo} \approx 0.2 m_{bo} c^2 \Gamma_j \approx 4 \times 10^{45} \left( \frac{m_{tail}}{4 \times 10^{-9} M_{\odot}} \right)^{1/2} \text{ erg}.$$ (10)

The observed emission implies $m_{tail} \sim 4 \times 10^{-7} M_{\odot}$.

Here we find that a fast ejecta tail is needed for a shock breakout to explain the observed γ-rays. However, the existence of such tail was suggested in the past based on theoretical models of the merger. A fast tail of the dynamical ejecta with $v_j \gtrsim 0.6 c$ is likely to arise from the interface of the merging neutron stars when the shock at the interface breaks out from the surface of the merging object (Kyutoku et al. 2014). Although it is hard to resolve numerically a small amount of fast moving components, some numerical simulations suggest that such a fast tail exists (Bauswein et al. 2013; Hotokezaka et al. 2013) and that it can contain as much as much as $0.15 \text{ cm}^2/\text{g}$.  

2 We find that by equating the diffusion time through the mass $m_{bo}$ with the dynamical time in the gas frame $R/c \Gamma_j$, assuming that $m_{bo}$ is spread over a width $R_{bo}/\Gamma_j$, $m_{tail} \approx 0.2 m_{bo} c^2$. 

3 While this paper was refereed Hotokezaka et al. (2018) analysed the highest resolution numerical simulations available of neutron star mergers (Kiuchi et al. 2017), finding a fast tail component that is similar to the one we use here.
as \( \sim 10^{-5}M_\odot \). Note that this amount depends on the fate of the central object after the merger and on the neutron star equation of state, e.g., more compact neutron stars eject more fast components and the amount of the shocked ejecta is significantly reduced when the merging neutron stars immediately collapse to a black hole (Hotokezaka et al. 2013).

Finally, we can estimate using these parameters what will be the characteristic temperature at the beginning of the spherical phase, namely, during the softer emission that follows the initial pulse. The rest frame density behind the shocked gas is roughly \( \rho_s \sim m_{\text{neut}} r_s^3 / [4 \pi R_0^3] \approx 2 \times 10^{-3} \text{gr cm}^{-3} \). The pressure, assuming \( \beta_p \approx 0.5 \) is \( p_s \approx 0.05 \rho_s c^2 \approx 10^{-3} \text{erg cm}^{-3} \) and the diffusion time, as measured in the gas rest frame, at the beginning of the spherical phase is \( \approx R_{\text{diff}} / c \tau_s \approx 2.5 \text{ s} \). Plugging these values to equation \( 5 \), we obtain a rest frame temperature of \( \sim 7 \text{ keV} \) and an observer frame temperature of about 7 keV.

4 NUMERICAL SIMULATIONS

The order of magnitude estimates, described earlier, demonstrate the potential of the model. We turn now to numerical simulations in order to provide a quantitative example which can be compared with the observations. We begin with relativistic hydrodynamic simulations. Then we post process the hydrodynamic results to obtain the observed \( \gamma \)-rays.

4.1 The relativistic hydrodynamic simulations

We have carried out relativistic hydrodynamic numerical simulations of jet propagation and cocoon formation within the ejecta. We focused on a choked jet since it has the potential to explain also the blue optical signal seen during the first day (Kasliwal et al. 2017). These simulations have used 2D axisymmetric geometry. We also carried out a single simulation of a successful jet to verify that it can also produce a significant \( \gamma \)-ray signal from the breakou t of its cocoon. The early phase of this simulation required 3D.

We used the public code PLUTO (Mignone et al. 2007), with an HLL Riemann solver and a third order Runge Kutta time stepping. Throughout the simulations we apply an equation of state with a constant adiabatic index of 4/3, as appropriate for a radiation dominated gas. We neglect gravity, as the gravitational dynamical times are longer than the typical interaction timescales.

4.2 The \( \gamma \)-ray emission

We use the results of the hydrodynamical simulations to calculate the shock breakout emission, assuming that diffusion of photons is radial, namely using the quasi-spherical symmetry approximation for photons diffusion. Following the shock breakout we find at each time step and each angle with respect to the jet axis, \( \theta \), the location within the flow from which photons can diffuse to the observer, which we term the diffusion depth. The photons escape to infinity (i.e., after crossing the diffusion depth) from the location where \( \tau = 1 \). For each time step and from each angle we emit the photons from the region that crossed the diffusion depth in the last time step (along the same angle) to the observer. The observed emission is found with a proper account for the light travel time and the Lorentz boost, according to the conditions at \( \tau = 1 \). To find the diffusion depth as a function of time we first identify the lab frame time of the shock breakout at this angle, \( t_{\text{bo}}(\theta) \). Upon the breakout the internal energy that is in the breakout layer (\( \tau = c / \gamma c_s \)) is released to the observer. After the breakout we find the diffusion depth by equating the time since the breakout to the diffusion time (see Appendix for details). To calculate the spectrum which is radiated from a given diffusion depth we use equation 5 and find \( T \), where the time for photon production is the time since the shock crossed the mass element at the diffusion depth and the hydrodynamical parameters are those that are at the diffusion depth at the time the photons are released. If the temperature obtained by this equation is higher than 50 keV we set \( T = 50 \text{ keV} \). The emitted spectrum is assumed to be a Wien spectrum, as expected before the radiation archive thermal equilibrium. The observed spectrum at any observer time, however, is obtained by integrating over the emission from the entire equal arrival time surface and is therefore not necessarily a Wien spectrum.

5 CHOKE D JET

Kasliwal et al. (2017) carried out numerical simulations which have shown that the interaction of a choked jet with the ejecta can simultaneously contribute to optical emission seen during the first day and generate a shock breakout that produces a \( \gamma \)-ray signal. Kasliwal et al. (2017) did not calculate the \( \gamma \)-ray emission directly form the simulations, instead they have shown that the breakout parameters are consistent with those that the order of magnitude estimates predict to produce a \( \gamma \)-ray signal with roughly the same energy, duration, hardness and delay as the one observed in GRB 170817A. Here we add a calculation of the \( \gamma \)-rays directly from the simulation (following the method explained above) in order to verify the order of magnitude estimates, including the predicted hard to soft evolution. The goal of this study is not to carry out a thorough scanning of the entire phase space in a search for an exact fit to the observed signal. Instead, we scan a relatively narrow part of the phase space to find if light curves with general properties that are similar to those observed can be generated.

The observed signal depends on the mass and velocity distributions of the ejecta as well as on the jet luminosity, opening angle and duration. It also depends on the delay between the merger time, at which the ejecta starts expanding, and the jet launch. The optical emission (and theoretical predictions) shows that the bulk of the mass is moving at 0.1-0.2c. The order of magnitude estimates require that a very small fraction of the fast ejecta tail mass would extend up to \( v_{\text{ej}, \text{max}} \approx 0.6-0.8c \), and similar amounts of mass are found to move at these velocities by some theoretical models (Bauswein et al. 2013; Hotokezaka et al. 2013; Kyutoku et al. 2014; Hotokezaka et al. 2018). Therefore, we consider an ejecta density profile that is composed of two components, a slow core in which \( M_c \) containing a few percent of \( M_{\text{ej}} \), and a fast low-mass tail in which \( M_f \) containing a few percent of \( M_f \). The optical emission during the first day depends on the delay of the jet launch as well as the jet and core ejecta properties, but is independent of the structure of the fast ejecta tail. The \( \gamma \)-ray emission depends on all the parameters including of course the tail. In order to consider configurations that can also account for the early optical emission we use similar jet properties and the same core structure as in Kasliwal et al. 2017 and vary only the tail structure. We verify that indeed the obtained optical/IR emission (using the same calculation method as in Kasliwal et al. 2017) is consistent with the observed one. Below, for completeness, we describe the full configuration we simulated.

Initially at \( t = 0 \) (defined as the merger time) we have a cold ejecta that expands radially. It is present from the base of the grid at \( r_{\text{esc}} = 4 \times 10^8 \text{ cm} \) up to \( r_{\text{max}} = 5.2 \times 10^9 \text{ cm} \). The ejecta has an
angular profile, where most of the mass (75%) is near the equator at \( \theta > 1.0 \) rad, where \( \theta \) is the angle with respect to the axis. The ejecta is divided also in the radial direction into two regions (with the same angular profile) - the main massive slow part that extends at \( t = 0 \) up to \( r_e = 1.3 \times 10^9 \) cm and a low-mass fast tail that extends at \( t = 0 \) between \( r_e \) and \( 5.2 \times 10^9 \) cm. The density profile of the dense part is:

\[
\rho(r, \theta) = \rho_0 r^{-2} \left( \frac{1}{4} + \sin^2 \theta \right),
\]

where \( \rho_0 \) is the normalization which is chosen for a total ejecta mass \( M_e = 0.1 M_\odot \). The velocity profile of the core is

\[
v_c(r) = v_{c, \text{max}} \frac{r}{r_e},
\]

where \( v_{c, \text{max}} = 0.2c \) is the maximal velocity of the core. The fast tail density profile has a very steep power-law in \( v \) between \( v_{c, \text{max}} \) and \( v_{c, \text{max}} \), and its normalisation is chosen so its total mass is \( M_e \).

Where needed we add an exponential (in density) transition layer between the core and the tail in order to have a continuous density profile. The jet is injected into the ejecta with a delay of 0.8 s for a total working time of 1 s and a total luminosity of \( L_j = 2.6 \times 10^{51} \) erg s\(^{-1} \). The jet is injected with a specific enthalpy of 20 at an opening angle of 0.7 rad from a nozzle at the base of the grid with a size of \( 10^6 \) cm.

We improve the resolution of the simulation in Kasliwal et al. (2017) as follows. In the \( r \)-axis we use 3 patches, the innermost one in the \( r \)-axis resolves the jet’s nozzle with 20 uniform cells from \( r = 0 \) to \( r = 2 \times 10^4 \) cm. The next patch stretches logarithmically from \( r = 2 \times 10^4 \) cm to \( r = 2 \times 10^{10} \) cm with 800 cells, and the last patch has 1200 uniform cells to \( r = 1.2 \times 10^{12} \) cm. In the \( z \)-axis we employ two uniform patches, one from \( z_{\text{beg}} = 4.5 \times 10^6 \) cm to \( z = 2 \times 10^{10} \) cm with 800 cells, and the second to \( z = 1.2 \times 10^{12} \) cm with 1200 cells. In total the grid contains \( 2020 \times 2000 \) cells, and the simulation lasts 40 seconds.

5.1 Hydrodynamics

At \( t = 0.8 \) s a jet is launched into the expanding ejecta, the jet is wide and covering a solid angle of about 25% of the entire sphere. A large fraction of the shocked material accumulates on top of the jet head and cannot be evacuated as it is not in a causal contact with the jet outer envelope (see top panel in figure 1). The wide jet is not collimated, propagating roughly conically inside the core as it shocks a significant fraction of it. After a total working time of 1 s the engine is turned off and within 0.5 s the jet is choked just before it emerges from the core ejecta depositing all the jet’s energy into the cocoon. The cocoon then breaks out of the core into the low-mass tail. No emission is released yet to the observer because to the high optical depth of the tail, but due to its low density the cocoon expands sideways and accelerates into the tail, in a way that is almost similar to expansion in a vacuum. First light is emitted upon the breakout of the cocoon from the fast ejecta tail (see bottom panel in figure 1). In the specific simulation depicted in figure 1 the shock breakout at \( \theta = 0.7 \) takes place at \( t = 6.2 \) s at a radius of \( 1.3 \times 10^{11} \) cm, corresponds to an observer time of \( \sim 1.8 \) s after the merger. At this point the shock is quasi-spherical and normal to the surface, crossing most angles at similar times, leaving only a fraction of unshocked ejecta around the equator. The velocity of the gas right behind the shock upon breakout is \( \Gamma \approx 2.0 \), but soon after the breakout it accelerates to \( \Gamma \approx 3.5 \).

5.2 \( \gamma \) -rays

Turning now to our main results we consider the \( \gamma \)-ray emission of the cocoon’s shock breakout. As mentioned earlier this emission depends on all the parameters including those of the faster tail that surround the main ejecta. We kept the jet and core parameters constant and checked the effect of the tail by considering several configurations (without doing an exhaustive parameter phase space search). We examined tail parameters in the following ranges: the density power-law \( -5 \) to \( -15 \), total mass \( (10^{-4} \text{ -- } 5 \times 10^{-2}) M_\odot \), and maximal velocity \((0.5 \text{ -- } 0.85)c\).

The outcome depends only on the parameters near the shock upon breakout, which are determined by these initial conditions. The light curves we obtained showed a large range of observed values, yet almost all light curves showed the expected common features of low-luminosity (compared to the total ejecta energy), low variability and hard to soft evolution. For the range of parameters we considered we find a large variation in the luminosity, where the peak luminosity varies between \( 10^{46} \) erg s\(^{-1} \) and \( 10^{49} \) erg s\(^{-1} \). Most simulations have shown hard to soft evolution with two spectral components. The ratio between the peaks of the two component is typically a few and varies between simulations by about an order of magnitude. The peak energy of the hard component is typically a few hundred keV, but in extreme cases it exceeds 1 MeV. The soft
component is typically lower than 100 keV but it may go under 1 keV in extreme cases. Smaller variations are seen in the duration and the delay, where the observed duration varies between 0.5 s and 4 s and the delay with respect to the merger between 1.5 and 4 s. The shape of the light curve also varies, Most have a fast rise and slow decay, but some have the opposite behavior and some are symmetrical. Some curves showed more structure than others but none have shown a rapid variability.

We find strong dependencies between the tail profile and the produced signal. First, a steeper tail density profile leads to a stronger shock, which in turn produces a brighter and harder signal. Another important parameter is the tail’s mass as the more massive ones stall the shock, filters less energetic material in high Lorentz factors, and result in a dimmer, later and longer signal. The tail’s front velocity has several effects. One is the time and duration of the peak as the shock would experience an earlier breakout with slower ejecta velocities. However, in these cases, where the shock breakout takes place in relatively small radii, the shock is more oblique and hence weaker. Additionally the shape of the signal varies greatly between the simulations, a lower ejecta front velocity mostly produces a slower rise and a sharp decline, while the fastest ejecta encounters a more spherical shock so that the angular time is shorter, the rise is sharp and the decay is gradual. Finally, in most, but not all, simulations we have carried out, the spherical phase dominates over the planar, both in time and luminosity, giving rise to a stronger soft component in the spectrum.

Among the configurations that we have examined there were many light curves that showed characteristics, such as duration, luminosity and hard to soft spectral evolution, to within an order of magnitude compared to those observed in GRB 170817A. In figure 2 we show an example of the light curve and spectrum observed at 0.7 rad that agree exceptionally well with all the observed properties of GRB 170817A. The resulting light curve starts rising at 1.5 s after the merger. It peaks at 2 s after the merger at $1.8 \times 10^{57}$ erg s$^{-1}$, and its total duration is $T_{\text{DD}} = 1.75$ s. It is composed of a bright initial pulse that lasts less than a second. It rises rapidly as the quasi-normal shock breaks out at $\theta \approx 0.6$. After the peak it drops gradually due to the combination of the spherical phase and shock breakout at $\theta \approx 0.8$, before a short sharp drop following the end of the spherical phase at these angles, after which a more gradual component, originated from wider angles, is observed for a bit longer than a second. The spectrum of the initial pulse is harder with $\nu L_\nu$ peaking at 110 keV, where the tail peaks at 60 keV, which corresponds to $T = 15$ keV. We stress that the agreement that we find does not imply that the setup which has been used in this simulation is the one we expect that took place in GRB 170817A. Since we did not scan the entire parameter space systematically, we expect that there are many other setups which are able to produce a similar or even better agreement with the observations.

5.3 UV/Optical/IR and Radio

As discussed in Kasliwal et al. (2017) the expanding ejecta produces a macronova signal powered by its radioactive decay, while its later interaction with the surrounding matter produces a radio afterglow (Hallinan et al. 2017). The macronova signal is boosted during the first day due to the mildly relativistic motion of the cocoon matter. The resulting optical/IR bolometric light curve and temperature, calculated from the same setup that produces the $\gamma$-rays depicted on figure 2 using the same methods described in Kasliwal et al. (2017), are shown in Fig. 3. Note that since we kept the same properties for the core of the ejecta as in Kasliwal et al. (2017) and similar jet properties the optical emission from all our choked jet simulations results in Opt/IR light curve that is in general agreement with the observations.

The interaction of the significant mildly relativistic outflow with the circum-merger material produces a strong radio afterglow signal (Nakar & Piran 2011; Piran et al. 2013; Hotokezaka & Piran 2015; Hallinan et al. 2017). This signal rises early because of the relativistic motion of the cocoon’s boosted material. Fig. 4 depicts the radio signal at 3 GHz, calculated from the same setup that produces the $\gamma$-ray signal depicted on figure 2, and compares it with the observations from Hallinan et al. (2017). The radio light curve is calculated using the same method as in Hallinan et al. (2017). The energy distribution as function of velocity of the outflow (i.e., $E(> \nu)$) is approximated to be spherically symmetric and is taken from the final snapshot of the simulation in the following way. For each Lorentz factor $\Gamma$ we measure all the energy that is within an opening angle of $1/\Gamma$ with respect to the line of sight and take its isotropic equivalent. Note that we used the profile of the expanding material obtained from a simulation that was chosen based on its $\gamma$-ray signal without any fit to the radio. The only free parameter we used in fitting the radio is the circum-merger density. In general the radio signal is very sensitive to the exact velocity profile of the ejecta and each simulation results in a different prediction for the radio. Some are in better agreement with the observations and some in worse. An interesting prediction of the choked jet model, which are seen in all our simulations, is that the signal continuous to rise in the near future.
6 SUCCESSFUL JET

We turn now to our simulation of a successful jet. Gottlieb et al. (2017) have shown that the cocoon of a successful jet also generates a mildly relativistic outflow at a wide angle\(^4\). They show that following the breakout the cocoon accelerates for a short time and after doubling its radius it sets into a homologous expansion. The cocoon cools adiabatically losing most of its internal energy and if there is no external source to dissipate the energy of the cocoon the remaining energy is released once the optical depth drops at frequencies far below \(\gamma\)-rays. We examine here a configuration that is similar to the one explored in Gottlieb et al. (2017) to which a fast low-mass tail has been added. We show that such a fast low-mass tail has a minor influence on the propagation of the cocoon once it emerges from the core of the ejecta. However, like in the case of a choked jet the breakout of the shock formed by the cocoon from the low-mass fast tail produces a \(\gamma\)-ray signal comparable to the observed \(\gamma\)-rays from GRB 170817A.

Gottlieb et al. (2017) have shown that jet propagation in the core of the ejecta must be done in 3D. We have used therefore simulation \(B\) of Gottlieb et al. (2017) up to the time the jet breaks out of the core. Since following the evolution to the radii that are needed in order to explore the shock breakout from the fast tail is too demanding for our 3D simulation we map the system at the time the jet breaks out of the core from 3D to 2D and continue simulating the jet propagation within the extended fast tail in 2D. We verify that this does not strongly affect the results (see below) Unlike the choked jet case we did not attempt to find a configuration that produces \(\gamma\)-rays with characteristics that are similar to those observed in GRB 170817A. Instead, we performed only a single simulation which shows that also a cocoon breakout from a successful jet generates a \(\sim 1\) s \(\gamma\)-ray signal with the correct delay which is bright enough and shows the hard to soft spectral evolution. We note, however, that this scenario has difficulties explaining the early UV/Optical/IR signal as the cocoon’s macronova in this case decays after several hours (see also Gottlieb et al. 2017) and it is not powerful enough to account for the bright UV/blue signal observed one day after the burst.

6.1 Initial conditions and numerical setup

The simulation is composed of two parts, the first, which follows the system up to the point the jet breaks out of the core ejecta is simulation \(B\) of Gottlieb et al. (2017), which is in 3D. The second is its later evolution in 2D. For completeness we provide here the initial configuration of simulation \(B\) of Gottlieb et al. (2017).\(^5\) At \(t = 0\), the time of the merger, the merger is surrounded by a cold ejecta starting from \(r_{\text{inf}} = 1.3 \times 10^8\) cm. The total mass of the ejecta is \(0.05 M_\odot\), with \(\rho \propto r^{-3.5}\). The ejecta outer surface is located at \(r = 3.9 \times 10^8\) cm with a velocity of \(0.2 c\). At \(t = 0.72\) s a narrow jet with an opening angle of \(10^\circ\) is injected into the system. The jet which has a specific enthalpy of 20 and total luminosity \(6.7 \times 10^{50}\text{erg s}^{-1}\), reaches the ejecta surface at \(9 \times 10^9\) cm after another 0.72 s. From this time onwards, the jet evolution is insensitive to the system’s dimensionality, as the jet has evacuated all the bulk mass.

\(^4\) Recently Lazzati et al. (2017b) also found a similar component, however in their model they do not identify an energy dissipation mechanism that will produce the observed \(\gamma\)-rays.

\(^5\) We use the freedom to scale numerical simulations (see Granot 2012) to scale the mass and length by five and three times the values used in Gottlieb et al. (2017), respectively.
in front of it so that the 2D numerical artifact of the “plug” (see discussion in Gottlieb et al. 2017) will not be present. We verify this in § B. We therefore utilize the snapshot of the 3D simulation at the time the jet breaks out of the core as our initial conditions for the 2D simulation. We convert the 3D results into 2D by averaging over rings along the rotation axis. Additionally we add to this snapshot the light, \(2 \times 10^{-3} \text{M}_\odot\), tail ahead of the core. The tail’s density profile is a power-law with \(\rho \propto r^{-10}\) and its front velocity is 0.8c, keeping the homologous profile by extending up to 4 breakout radii.

The numerical setup (solver, equation of state etc.) is identical to the choked jet simulation. The grids are however somewhat different as to reflect the earlier 3D simulations. The grid is divided into three patches in each axis, while the first two are identical to the original 3D simulation. The innermost patches are distributed uniformly in \(r\) (50 cells) and \(z\) (400 cells) axes, extending to \(3 \times 10^8\) cm and \(6 \times 10^8\) cm, respectively. The \(z\)-axis begins at \(1.3 \times 10^9\) cm. The second patches are logarithmic with 240 and 600 cells up to \(9 \times 10^{10}\) cm and \(1.2 \times 10^{11}\) in \(r\) and \(z\) axes, respectively. The extension of the grid to include the ejecta tail is to \(1 \times 10^9\) cm.

6.2 Hydrodynamics

We have injected at 0.72 s after the merger, a narrow (\(\theta_j = 10^\circ\)) jet into the expanding ejecta (Fig. 5). The jet is well collimated and able to evacuate efficiently the ejecta in front of it and propagate at mildly relativistic velocities until breaking out of the core ejecta within another 0.72 s, before its engine is turned off 1 s after the launch. At this point the jet enters the dilute extended tail, and accelerates to a Lorentz factor of a couple of dozens. The jet is accompanied by a hot cocoon that expands to a wide angle and moves in mildly relativistic velocities. The cocoon shape is aspherical, and the shock breakout is oblique. It does not reach angles larger than \(\pi/4\). However it is fast, and its Lorentz factor is almost 3 upon breakout and 5 after the acceleration phase.

In Fig. 5 we show the breakout at \(\theta = 0.7\) rad. It takes place after 9.8 s and at \(r = 2.4 \times 10^{11}\) cm, corresponds to \(t_{\text{obs}} \approx 1.8\) s. The main differences from the choked jet case (Fig. 1) are the initial jet collimation, (shown in the top panel) and its presence in the homologous phase on the \(z\)-axis (shown in the bottom panel) at the time of the cocoon breakout with a width of slightly more than a light second, and the cocoon which is less spherical.

6.3 \(\gamma\)-rays

We calculate the \(\gamma\)-ray emission arising from a shock breakout of the cocoon from the extended tail at large angles, where the emission from the jet itself does not contribute at all. In Figure 6 we present the signal for an observer at \(\theta_{\text{obs}} = 0.7\) rad. The delay of slightly less than two seconds, and the duration \(T_{\gamma} = 1.6\) s are similar to GRB 170817A. The light curve shape in this simulation is determined mostly by the obliqueness of the shock. The fast rise to the peak is due to the shock at \(0.55 < \theta < 0.7\); the peak is maintained by angular contribution from the shock at \(0.4 < \theta < 0.55\), followed by a steep decline as the shock does not reach angles larger than \(\pi/4\). With a peak luminosity of \(9 \times 10^{52}\) erg s\(^{-1}\) the signal from the simulation is brighter by about an order of magnitude compared to GRB 170817A. The spectrum shows a clear hard to soft evolution, but both components are harder by an order of magnitude compare to GRB 170817A. Dividing the spectra at \(t = 2.3\) s, during the sharp drop from the peak of the signal, the hard component is about 1MeV, while the softer one is several hundreds of keV.

Given that (i) we have used an existing 3D simulation as our initial condition and (ii) we did not do any parameter search but we run only a single set of parameters for the extended tail, the fact that most features of GRB170817 are present and fit up to better than an order of magnitude with this model is exceptional. Scanning the parameters space carefully is most likely to yield a significantly better match with the observed \(\gamma\)-ray signal of GRB 170817A.

7 DISSOCIATION OF THE NUCLEI

Before concluding we note a possibility, that we did not consider in this work, which is the dissociation of heavy nuclei by the cocoon’s shock. The internal energy per baryon in the shocked ejecta is \(\gtrsim 10\) MeV. This exceeds the binding energy per baryon of heavy nuclides \(\sim 8\) MeV. In principle there is enough energy to dissociate heavy nuclei in the shocked ejecta. As most of the energy is stored in the photons, the condition of the photodissociation of heavy nuclei depends sensitively on the temperature. The reaction rate is proportional to the number of photons above \(\sim 8\) MeV, which is strongly suppressed in the exponential tail of the photon’s spectrum. At thermal equilibrium, which is expected at the shocked ejecta, the crit-
the ejecta in its own frame is of a wide jet) the velocity of the forward shock driven into
Assuming that the jet is not strongly collimated (as in case ical temperature, above which photodissociation occurs, is \( \sim 200 \) keV (e.g., Woosley & Howard 1978).

The jet propagates sub-relativistically within the core ejecta. Assuming that the jet is not strongly collimated (as in case of a wide jet) the velocity of the forward shock driven into the ejecta in its own frame is \( \beta' \sim \left[ L_{j,iso}/(4\pi R^2\rho c^5)\right]^{0.5} \sim 0.03L_{j,iso}^{1/2}(M_j/0.1M_\odot)^{-1/2}/(R/10^9\text{cm})^{1/2} \) where \( L_{j,iso} \) is the jet isotropic equivalent luminosity (Bromberg et al. 2011). The temperature behind the shock is estimated as \( T \sim (\rho_j v_j^2/\eta BB)^{1/4} \sim 150\rho_j^{1/4,v_j^{1/2}}(R/10^9\text{cm})^{-1/2}\text{keV} \). If the jet is collimated than it prop-
agrases faster and the temperature is higher. This implies, giving the high sensitivity on the temperature, that photodissociation is possible. In our simulations of choked and successful jets the delay is almost 1s so the jet interacts with the ejecta at radii > 10^9 cm and the typical temperature is \( \approx 100\text{keV} \), thus we do not expect photodissociation to take place. We stress, however, that as this process is sensitive to the temperature and density evolution, the final abundance of nuclei should be addressed with more detailed calculations.

In addition to photodissociation, free neutrons may also play an important role to disintegrate heavy nuclei. In the shocked material, free neutrons from the upstream have a velocity of \( \overline{v}/(0.1c) \). They thermalize through collision with heavy nuclei. During this thermalization process, heavy nuclei are disintegrated. This process occurs until a few hundreds ms after the mass ejection, by which free neutrons are likely exhausted for the typical ejecta parameters (\( v_{ej} < 0.5c \)). Note that free neutrons can remain in the fast components of the ejecta with \( v_{ej}\geq0.5c \) (Metzger et al. 2015). Energetic free neutrons in the shock disintegrate heavy nuclei within the fast tail even at later times.

To conclude the jet propagation may affect the ejecta composition. This effect is only along the jet path, so a wide choked jet may affect a larger portion than a narrow collimated jet. In any case most of the ejecta, mostly the part which propagates at low latitudes will not be affected. If the jet dissociates heavy nuclei in high latitude ejecta it will change its composition. The heavy nuclei will disappear and this will reduce the opacity and affect the radioactive heating rate within the region influenced by the cocoon. This will clearly affect the UV/optical/IR macronova signal especially at early times. Such process may be related to the blue light observed during the first day. The composition, however, will not strongly influence the \( \gamma \)-rays that arise from the shock breakout or the late radio emission which arises due to the interaction of the ejecta with the surrounding matter.

8 CONCLUSIONS

The short GRB 170817A was a not a regular sGRB. It was a ll-GRB, namely a low luminosity one. Like llsGRBs (low-luminosity long GRBs) that are not produced by the same mechanism as regular long GRBs, llsGRBs are not produced by the same mechanism as regular sGRBs. In fact we suggest that, while the astrophysical scenarios are very different (a Collapsar vs a merger) there is a similarity between the physical mechanism that produces the two types of low-luminosity GRBs. Both are produced by a shock breakout. In the former the shock breakout from the envelope of the star. In the latter the shock breakout is from the surrounding matter (ejec-
ta) that is thrown out to space during the merger process.

We have first shown, using the traditional compactness argument, that the observed \( \gamma \)-rays implied that the event involved at least mildly relativistic outflow. At the same time this was not a regular sGRB viewed from the side (Kasliwal et al. 2017). There is no simple way to obtain a spherical relativistic outflow in the configuration involved. To accelerate spherically the outer layers of the ejecta to \( \gamma\gtrsim2 \) we will need to accelerate the whole bulk of the ejecta to this velocity and this is inconsistent with the observations. Thus, we conclude that the system must have involved a relativistic jet that carries the energy through the bulk of the ejecta to the outer layers and deposits it there, accelerating only a very small fraction of the ejecta. This jet might have penetrated success-
fully the ejecta producing a sGRB that was pointing away from us (and not observed by us), or it might have been completely choked within the outflow. In either case the jet would have produced a hot cocoon within the ejecta and we suggest that the observed \( \gamma \)-rays arose during the breakout of this cocoon from the fast outer tail of the ejecta.

We have outlined here an order of magnitude model for a relativistic shock breakout from a moving ejecta, highlighting the differences between a shock that emerges from a static medium to one that emerges from an expanding one. We have shown that three generic properties of such a shock breakout are: (i) The light curve is smooth. It may show some structure but no fast variability; (ii) Only a small fraction of the total available energy is emitted at this stage; (iii) The emission involves two phases, a planar phase with a hard spectrum and a spherical phase that has a softer thermal spectrum. These feature resembles well the observed features of the observed llsGRB 170817A. It is important to note that the
shock breakout scenario has four parameters that control its emission: The Lorentz factor of the shock, $\Gamma_s$, the velocity of the ejecta, $\beta_{e,\text{max}}$, the radius in which the shock breakout takes place, $R_{\text{bo}}$, and the causally connected mass that is within $R_{\text{bo}}/\Gamma_s$. Moreover, there is little freedom in the amount of this mass allowed by reasonable ejecta models, and the observables (e.g. the total energy) are rather weakly dependent on it. These should be compared with five observables, the total energy, the duration, the delay after the merger, the peak energy of the hard component and the peak energy of the soft component. Thus, the model is over-constrained and the basic agreement of its predictions with the observations points out in its favor. Note that the model requires a fast moving low-mass tail surrounding the main ejecta. However, as pointed out in §3 such fast moving material is expected (Kytotoku et al. 2014; Hotokezaka et al. 2018) and was possibly noticed in some numerical simulations (Bauswein et al. 2013; Hotokezaka et al. 2013).

We then carried out numerical relativistic hydrodynamics simulations that follow the propagation of a jet through the expanding ejecta, the formation of a cocoon and the breakout of the shock produced by the cocoon from the ejecta. We post-processed the hydrodynamic simulation to calculate the resulting $\gamma$-ray emission. For the choked jet scenario we kept the parameters of the jet and the main component of the ejecta (its massive core) and have explored a small part of the phase space of fast tail configurations. Remarkably we have found over a large fraction of the phase space that we explored a $\gamma$-ray signal that is comparable to within an order of magnitude with the one observed in GRB 170817A.

The observed associated macronova (e.g. Kasliwal et al. 2017) was also somewhat different from the expectations. It had a strong early UV/Opt signal that was much brighter and bluer than earlier expectations. Kasliwal et al. (2017) showed that a cocoon from a choked jet can also account for this early blue light. This is mostly due to a relativistic boost given at early times by the fast cocoon material. This boost makes the macronova signal both brighter and bluer, before the emission of the slower ejecta dominates at later times. All our choked jet simulations here, which have similar jet and ejecta core parameters as those of Kasliwal et al. (2017), recover the observed macronova. The mildly relativistic cocoon material interacts later with the surrounding matter and produces a radio signal. This signal is sensitive to the velocity distribution of the cocoon, which in turn depends on all the jet and ejecta parameters. We calculated the radio emission predicted for one of our simulations, that fit both the $\gamma$-rays and the macronova, and found that it is also in general agreement with the observed radio emission (Hallinan et al. 2017). We find it remarkable that within a single set of parameters we find agreement with the observed $\gamma$-rays, macronova and radio signals. A clear prediction of the choked jet scenario is that the radio signal will keep increasing over the near future.

While we have focused on the case of a choked jet, due to its ability to explain also the macronova, we have also shown that a rather similar $\gamma$-ray signal can be generated by a cocoon shock breakout driven by a jet that emerges and produces a sGRB pointing away from us. Here we have studied only one configuration whose $\gamma$-ray features resemble the observations. A successful jet produces a much less massive and energetic cocoon and is therefore not expected to affect the macronova emission for more than several hours after the merger. Its lower energy also predicts a fainter radio signal which may be dominated by the off-axis radio afterglow of the relativistic jet.

We also examined briefly the effect of the jet propagation on the ejecta composition. We find that the temperature in the shocked core ejecta is marginal for photodissociation of the heavy nuclei. A luminous jet ($L_{\nu,100} \gtrsim 10^{52}$ erg/s) that is launched quickly after the merger (within $\sim 0.1$ s) is expected to dissociate a significant fraction of the ejecta nuclei that lie in its path, while a less luminous jet and/or a longer delay will most likely have a minor effect on the ejecta composition. In the scenarios we considered here the jet was launched about 1s after the merger and thus the temperatures found in our simulations are too low for photodissociation. It is possible, however, that for other events, or even for this one, the conditions are such that photodissociation takes place. If it does then it may strongly affect the early optical emission (which may be related to the observed first day blue light), but not the late optical/IR emission nor the $\gamma$-rays or the radio.

Finally we note that while the cocoon breakout $\gamma$-ray signal is much wider than the emission of a regular sGRB (in both cases of choked and successful jets), it is not isotropic. The signal depends on the observer angle and nearer to the jet’s axis it is typically brighter, harder, shorter and with a smaller delay (clearly in the case of a successful jet the sGRB emission from the jet dominates completely over this emission for an on-axis observer). It is dimmer for larger observing angles (a detailed study of the emission as a function of the viewing angle will be published elsewhere), yet for some parameters it may be observed over the entire $4\pi$. Overall we expect that the $\nu$sGRB signal will be observed over a quite large solid angle and will accompany many GW events.

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APPENDIX A: THE $\gamma$-RAY SHOCK BREAKOUT EMISSION

To calculate the $\gamma$-ray emission from the shock breakout we begin by monitoring the shock location as a function of time and angle. Then, for each angle $\theta$ we find the shock breakout time $t_{bo}(\theta)$ and radius $r_{bo}(\theta)$. For each lab frame time $t > t_{bo}(\theta)$, the emission from angle $\theta$ originates in $r_{e}(\theta)$, the radius from where photons diffuse out to the photosphere ($r_{ph}(\theta) = r(\theta$'s time-travel time from the photosphere. We assume that the photons diffusion to be radial. Hence the location $r_{e}(\theta)$ is determined by

$$r_{e}(\theta) = r_{ph}(\theta) - \frac{t - t_{bo}(\theta) c}{\Gamma(r_{e})}, \quad (A1)$$

where $\Gamma(r_{c})$ is the Lorentz factor of the emitting region. The released rest-frame energy per solid angle $d\Omega$ between two successive time-steps $t_{1}$,$t_{2}$ is the total thermal energy $4\rho$ of the emitting region:

$$\frac{dE}{d\Omega}(t_{1}, t_{2}, \theta) = \int_{r_{e}(t_{1})}^{r_{e}(t_{2})} \frac{4 \rho(r) \Gamma(r)^{2}}{r^{2}} dr, \quad (A2)$$

where $r_{e}(t_{bo}(\theta)) = \infty$. The observed energy is then obtained by boosting the rest-frame energy $dE$ to the observer, and the arrival time is given by considering the light-travel time from the photosphere. The observed bolometric luminosity is then calculated by integrating over all times and angles.

APPENDIX B: CONVERGENCE TEST

Computer time limitations make it impossible to carry the full simulations in 3D. Therefore, in the collimated jet simulation, where 3D are necessary as long as the jet is collimated, we switch from 3D to 2D when the configuration becomes insensitive to 3D effects, that is once a successful jet breaks out from the ejecta. We verify that converting the 3D simulation of the successful jet to 2D upon breakout from the core ejecta does not heavily affect the simulation outcome by performing two tests. First, we verified that a 2D simulation with a low-mass tail and one without one give similar results in their energy distributions at different angles. Then, we run identical 2D and 3D simulations of the jet in a setup where there is no fast tail to the ejecta, starting at the time the jet and the cocoon break out of the core and until the jet increases its radius by a factor of 10 (initial conditions are taken from the 3D simulation). In figure B1 we compare the 3D and 2D velocity and energy distributions at various angles and found them to have a high degree of similarity at the last snapshots of the two simulations. The biggest discrepancy is found in small velocities, around $\Gamma_{\beta} = 0.1$, where the material has not reached the homologous phase yet. This component does not contribute to the $\gamma$-ray emission.