Imprints of r-process heating on fall-back accretion: distinguishing black hole–neutron star from double neutron star mergers

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ABSTRACT

Mergers of compact binaries containing two neutron stars (NS–NS), or a neutron star and a stellar mass black hole (NS–BH), are likely progenitors of short-duration gamma-ray bursts (SGRBs). A fraction $\gtrsim 20$ per cent of SGRBs is followed by temporally extended ($\gtrsim$ minute-long), variable X-ray emission, attributed to ongoing activity of the central engine. One source of late-time engine activity is fall-back accretion of bound tidal ejecta; however, observed extended emission light curves do not track the naively anticipated, uninterrupted $t^{-5/3}$ power-law decay, instead showing a lull or gap in emission typically lasting tens of seconds after the burst. Here, we re-examine the impact of heating due to rapid neutron capture ($r$-process) nucleosynthesis on the rate of the fall-back accretion, using ejecta properties extracted from numerical relativity simulations of NS–BH mergers. Heating by the $r$-process has its greatest impact on marginally bound matter, hence its relevance to late-time fall-back. Depending on the electron fraction of the ejecta and the mass of the remnant black hole, $r$-process heating can imprint a range of fall-back behaviour, ranging from temporal gaps of up to tens of seconds to complete late-time cut-off in the accretion rate. This behaviour is robust to realistic variations in the nuclear heating experienced by different parts of the ejecta. Central black holes with masses $\lesssim 3M_\odot$ typically experience absolute cut-offs in the fall-back rate, while more massive $\gtrsim 6-8M_\odot$ black holes instead show temporal gaps. We thus propose that SGRBs showing extended X-ray emission arise from NS–BH, rather than NS–NS, mergers. Our model implies an NS–BH merger detection rate by LIGO that, in steady state, is comparable to or greater than that of NS–NS mergers.

Key words: accretion, accretion discs – hydrodynamics – nuclear reactions, nucleosynthesis, abundances – gamma-ray burst: general.

1 INTRODUCTION

Short-duration gamma-ray bursts (SGRBs) are commonly believed to be powered by rapid accretion on to a rapidly spinning black hole, following the coalescence of a compact binary system (e.g. Narayan, Paczynski & Piran 1992). The latter may be comprised either of two neutron stars (NS–NS) or a neutron star and a stellar mass black hole (NS–BH). The recent discovery of an SGRB (Abbott et al. 2017b) coincident with the gravitational wave event GW170817 (Abbott et al. 2017a), as well as non-thermal emission from the off-axis afterglow of a relativistic jet (e.g. Alexander et al. 2017; Haggard et al. 2017; Hallinan et al., in preparation; Margutti et al. 2017; Troja et al. 2017), provided compelling evidence that at least some SGRBs arise from NS–NS mergers (Blinnikov et al. 1984; Goodman 1986; Paczynski 1986; Eichler et al. 1989). Although no NS–BH binaries are currently known, they are theoretically predicted to exist and may also give rise to SGRB emission in cases where the merger results in the creation of an accretion disc (e.g. Rosswog 2005; Kyutoku et al. 2011; Foucart et al. 2013, 2014, 2016; Paschalidis, Ruiz & Shapiro 2015; Bhattacharya, Kumar & Smoot 2018). This latter condition requires that the BH be of sufficiently low mass and/or rapidly spinning in the prograde direction with respect to the binary orbit, such that the NS will be tidally disrupted before falling into the BH horizon (e.g. Foucart 2012).

At least $\sim 20$ per cent of SGRBs are followed by temporally extended X-ray emission, which lasts $\sim 10$–1000 s or longer after the initial prompt gamma-ray burst (Gehrels et al. 2006; Norris & Bonnell 2006; Perley et al. 2009; Minaev, Pozanenko & Loznikov 2010; Norris, Gehrels & Scargle 2010; Kaneko et al. 2015; Kisaka, Ioka & Sakamoto 2017; Burns et al. 2018); Fig. 1 shows the light curves in two clear cases. This ‘extended emission’ (EE) is too...
et al. 2009) and GRB 060614 (Gehrels et al. 2006), which show their object remnant left by the merger (however, see Eichler, Guetta & Metzger, Thompson & Quataert 2008; Bucciantini et al. 2012; Gao et al. 2013; Rowlinson et al. 2013; Gompertz, O’Brien & Wynn 2016). A stable magnetar is also disfavoured in place constraints on the amount of rotational energy released from the lack of late-time radio detections of short GRBs is beginning to indicative of black hole formation (Metzger & Fernández 2014b); however, at earlier times, less than a few seconds following mass ejection from the merger, the r-process heating rate is orders of magnitude higher. Energy released by the r-process does not qualitatively alter the dynamics of the bulk of the unbound ejecta (Rosswog et al. 2014). However, it can increase the quantity of ejecta and critically shape the dynamics of the marginally bound ejecta responsible for fall-back accretion, particularly on the second to minute timescales of relevance to the observed EE (Metzger et al. 2010a).

Depending on the electron fraction of the ejecta, $Y_e \sim 0.02–0.3$, the r-process releases a total energy of $Q_{\text{tot}} \sim 1–3\text{MeV}$ per nucleon, mostly through beta-decays, at an approximately constant rate over a characteristic heating time-scale $t_{\text{heat}} \sim 1\text{s}$ following ejection (see Fig. 2). Then, at times $t \gg t_{\text{heat}}$ (e.g. relevant to the kilonova), the heating rate approaches an asymptotic power-law decay $\dot{q} \propto t^{-1.3}$ (Metzger et al. 2010b). A nucleon (mass $m_n$) that is marginally gravitationally bound to the black hole of mass $M$, on an orbit of energy per nucleon $|E_{\text{isol}}| = GMm_n/2a$ and semimajor axis $a$ (under the approximation of Newtonian gravity), returns to the hole and circularizes into the accretion disc on a time-scale given by its orbital period:

$$t_{\text{heat}} = 2\pi \left( \frac{a^3}{GM} \right)^{1/2} \approx 1.6s \left( \frac{|E_{\text{isol}}|}{1\text{MeV}} \right)^{-3/2} \left( \frac{M}{5M_\odot} \right).$$

This expression reveals several important facts. First, the total energy available from the r-process, $Q_{\text{tot}}$, exceeds the binding energy of orbits with fall-back times comparable to observed EE after SGRBs ($\gtrsim 10\text{s}$; Fig. 1); including the effects of r-process heating is thus crucial to determining the late-time fall-back rate (Metzger et al. 2010a). Also of interest is the apparent coincidence that the fall-back time of matter with energy $|E_{\text{isol}}| \sim Q_{\text{tot}}$ is comparable to the time-scale $t_{\text{heat}} \sim 1\text{s}$ over which the bulk of the heating occurs. This means that different parts of the debris could receive different amounts of the total available heating, depending on their fall-back time. Metzger et al. (2010a) show that this can imprint a more complex mass fall-back evolution than the standard $\propto t^{-5/3}$ decay, instead generating either temporal "gaps" of several seconds or sharp cut-offs in the fall-back rate after a certain time, depending on the ratio of $t_{\text{heat}}(|E_{\text{isol}}| = Q_{\text{tot}})$ and $t_{\text{heat}}$.

1We use the term gap throughout this paper to describe the epoch between the merger and the light-curve maximum, when the X-ray luminosity, while non-zero, is substantially suppressed below the naive power-law expectation.
In this paper, we apply the model of Metzger et al. (2010a) in order to estimate the effects of r-process heating on the energy distribution of the merger ejecta and its resulting mass fall-back rate, using for the first time initial conditions for the debris properties taken directly from a numerical relativity simulation of an NS–BH merger (Foucart et al. 2016). We also explore what effects the Y_e-dependent nuclear heating rate has, given a realistic spread in the debris properties, on the predicted range of fall-back behaviour.

In Section 2, we describe the problem set-up, our treatment of the nuclear heating, and the numerical technique. In Section 3, we present our results for the mass fall-back rate. Finally, in Section 4 we map our findings on to mergers leaving BHs of different mass scales to demonstrate how NS–NS and NS–BH mergers might in principle be distinguished based on the properties of their late-time X-ray light curves.

2 MODEL

2.1 Numerical simulation data

We use 3D position/velocity data taken from the grid points of the NS–BH merger simulation ‘M5-S7-I60’ performed by Foucart et al. (2014) using the numerical code SPEC (Kidder et al. 2000). This simulation implemented the DD2 equation of state (Hempel et al. 2012) for a 1.4 M⊙ NS, and included a neutrino leakage scheme, as implemented in Deaton et al. (2013). The initial mass of the BH is M_BH = 5 M⊙. This is a precessing system, with a dimensionless spin of η = 0.7 on the BH, prograde but misaligned by 60° with respect to the orbital angular momentum. The final mass of the BH after the merger is M = 6.11 M⊙.

We extract data on the ejecta tidal tail at a time t = 15 ms after merger. These are then used as initial conditions in post-processing analysis, in which we evolve the fluid elements further as non-interacting Lagrangian particles. Because initial energies, positions, and velocities used in our analysis were derived from a general relativistic simulation, they are inconsistent with the Newtonian expression for total energy, We therefore rescale the initial velocities from the simulations (v_{GR}) to their Newtonian equivalent according to

$$v_N = \frac{v_{GR}}{\gamma_{GR}} |v_{GR}|$$

where

$$\gamma_{GR} = \sqrt{1 + \frac{2GM_{BH}}{r}}$$

and $E_{tot} = -m_b(u + 1)$ is an estimate of the specific binding energy (per nucleon) of the test particle, assuming geodesic motion in a time-independent spacetime, and r is the distance from the BH.

Fig. 3 shows the distribution of initial ejecta properties as a function of $E_{tot}$ and electron fraction Y_e, weighted by mass. Most of the debris is neutron-rich, with an electron fraction $Y_e \lesssim 0.05$. A dashed line separates ejecta that is initially gravitationally bound ($E_{tot} \geq 0$) from unbound material ($E_{tot} < 0$). Of the total mass 0.116 M⊙ of the tidal ejecta, approximately 1.23 × 10^{-2} M⊙ is

2The Spectral Einstein Code: http://www.black-holes.org/SpEC.html.
unbound. Even if they are initially bound to the black hole, fluid elements above the solid blue line could in principle gain sufficient heating from the r-process (Fig. 2; bottom panel) during their orbit to become unbound. As we discuss later, depending on the time-scale over which the r-process heating is released, matter near this line can also remain bound but fall back to the black hole at a later time than had it experienced zero heating.

2.2 R-process heating

We include the effects of r-process heating on the trajectories of the fall-back debris in a post-processing step, similar to the model employed by Metzger et al. (2010a). The total nuclear energy released as neutrons are captured on to seed nuclei during the r-process is approximately given by the difference between the initial and final nuclear binding energies:

\[ Q_{\text{tot}} \simeq (1 - f_\gamma) \left[ \left( \frac{B}{A} \right)_s - X_s \left( \frac{B}{A} \right)_s - X_s \right] . \]  

(3)

minus the fraction \( f_\gamma \) of energy lost to neutrino emission (\( f_\gamma \approx 1/2 \) at early times of interest). Here, \( X_s = AY_e/Z \) is the mass fraction of seed nuclei of average atomic mass number \( A \) and charge \( Z \); \( (\frac{B}{A})_s = 8.7 \) and \( (\frac{B}{A})_s = 8 \) MeV nuc\(^{-1} \) are the average binding energies of seed and final r-process nuclei, respectively; \( X_s = 1 - X \) is the neutron mass fraction; and \( \Delta_a = (m_n - m_p)c^2 = 1.293 \) MeV is the neutron–proton mass difference. For typical values \( Y_e \approx 0.05 \) (Fig. 3), \( \tilde{A} \approx 90, \tilde{Z} = 38, f_\gamma = 0.45 \), we find \( Q_{\text{tot}} \approx 3 \) MeV, consistent with the results of SkyNet nuclear reaction calculations (Lippuner et al. 2017) shown in the bottom panel of Fig. 2.

As shown in the top panel Fig. 2, the r-process heating in freely expanding unbound debris is approximately constant for a time-scale of \( t_{\text{heat}} \sim 1 \) s, when most of the total energy is released, before rapidly entering a power-law decline. We approximate this behaviour by a function of the simple form:

\[ \dot{q} = \begin{cases} \frac{Q_{\text{tot}}}{t_{\text{heat}}}, & \text{if } t \leq t_{\text{heat}} \text{ and } v_t > 0, \\ 0, & \text{if } t > t_{\text{heat}} \text{ or } v_t < 0. \end{cases} \]

(4)

where \( v_t \) is the radial velocity of the fluid element and the values of \( Q_{\text{tot}} \sim 3 \) MeV and \( t_{\text{heat}} \sim 1 \) s are parameters that we allow to vary modestly about these fiducial values.

The r-process heating, as described by equation (4), is assumed to terminate if and once matter starts to return to the black hole (\( v_t < 0 \)), for reasons we now discuss. Although the heating rate evolution shown in Fig. 2 was calculated for unbound debris, this behaviour is a good approximation also for bound debris during its initial outward motion. However, the heating is abruptly suppressed once matter reaches apocenter and begins to return to the black hole (Metzger et al. 2010a). As matter undergoes re-compression, its temperature rises adiabatically and the r-process path (which is determined at a fixed Z by the equilibrium between neutron capture (\( n, \gamma \)) and photodissociation (\( \gamma, n \) processes) is driven closer to the stable valley, where the \( \beta \)-decay time-scales, and thus neutron consumption and energy release time-scale, becomes much longer. Metzger et al. (2010a) show that to good approximation the heating rate effectively shuts off once \( v_t < 0 \), motivating us to neglect heating entirely during re-compression.

2.3 Numerical model

At early times after the merger, the ejecta is dense and highly opaque to photons, such that all of the r-process heating (other than that which escapes as neutrinos) goes into internal thermal energy. The fluid element orbits of interest possess high eccentricities \( e \), where

\[ 1 - e = \frac{r_p}{a} \approx 0.03 \left( \frac{r_p}{3r_g} \right) \left( \frac{|E_{\text{tot}}|}{3 \text{ MeV}} \right) . \]

(5)

Here, we have normalized the pericenter radius of the debris, \( r_p \), to the gravitational radius of the BH, \( r_g = GM/c^2 \). Temperature and density gradients in the ejecta are thus directed nearly radially outwards, such that the r-process energy will be transferred through PdV work into ejecta kinetic energy on the local expansion time-scale. Fig. 4 shows a schematic illustration of the influence of r-process heating on the trajectories of different fluid elements.

Following explicitly the effects of r-process heating on the debris dynamics would require a 3D hydrodynamical simulation across a large dynamical range in radius. However, given the highly supersonic expansion velocities of the ejecta, such a treatment is not necessary, as the main effect of the heating is a local slow acceleration of the ejecta along the local pressure gradient radial direction and thus the adiabatic conversion of the injected thermal energy to kinetic energy. Starting with fluid element velocities rescaled from the simulation data (equation 2), we directly increase the kinetic energy of the \( i \)th element according to

\[ \frac{d}{dt} \left( \frac{1}{2} m_n v_{ti}^2 \right) = \dot{q}_i , \]

(6)

where \( \dot{q}_i \) follows equation (4). For each fluid element, we follow its 3D trajectory until it reaches either periapse, which we define as the fall-back time \( t_{fb} \), or once the simulation terminates. We then record \( t_{fb} \) for all bound material and use the total mass in different bins of \( t_{fb} \) to calculate the fall-back accretion rate.
fall-back history. In particular, there are two possible outcomes in shape of the fall-back curve: either an absolute cut-off after some time (e.g. as in the $Q_{\text{tot}} = 2\, \text{MeV}$, $t_{\text{heat}} = 2\, \text{s}$ case), or a cut-off followed by re-emergence of fall-back, i.e. a ‘gap’ (e.g. as in the $Q_{\text{tot}} = 1.5\, \text{MeV}$ and $t_{\text{heat}} = 1.5\, \text{s}$ case). As we now discuss, these qualitatively different behaviours can be understood by comparing the time-scale over which the ejecta is heated to its orbital time-scale (equation 1; see also Metzger et al. 2010a).

First, consider the existence of a critical orbital period $t_{\text{orb,c}}$, which corresponds to matter bound to the BH by an energy equal to the energy $\approx (Q_{\text{tot}}/t_{\text{heat}})\eta_{\text{orb,c}}$ it receives from the $r$-process over the orbital period $t_{\text{orb}}$ (when $t_{\text{orb}} \gtrsim t_{\text{heat}}$). Using equation (1) for $t_{\text{orb}}$, this gives

$$t_{\text{orb,c}} \approx 0.62\, \text{s} \left(\frac{Q_{\text{tot}}}{3\, \text{MeV}}\right)^{-3/5} \left(\frac{M}{5\, M_\odot}\right)^{2/5} \left(\frac{t_{\text{heat}}}{1\, \text{s}}\right)^{3/5}. \quad (7)$$

Ejecta that starts on an orbit of period $t_{\text{orb}} \gg t_{\text{orb,c}}$ always receives the full $r$-process heating before reaching apocenter (at which point the matter starts to re-compress and heating shuts off for reasons discussed earlier), while matter that starts very tightly bound ($t_{\text{orb}} \ll t_{\text{orb,c}}$) may receive only a fraction $t_{\text{orb}}/t_{\text{heat}}$ of the total heating (if $t_{\text{heat}} \gtrsim t_{\text{orb,c}}$).

Crucially, however, if $t_{\text{orb}} \gtrsim t_{\text{orb,c}}$ even matter with an initial $t_{\text{orb}}$ that is slightly less than $t_{\text{orb,c}}$ can also receive the full heating because as the energy of a fluid element increases, its orbital period also grows, giving it more time to receive the full allotment of nuclear energy (in other words, the final value of $t_{\text{orb}}$ diverges in a runaway process due to the $r$-process heating). This preferential heating opens a gap in the energy distribution of the debris, which can result in an absolute cut-off in the accretion rate, or a temporal gap, depending on the location of the gap.

Whether such behaviour is possible depends on whether the $r$-process heating indeed acts uniformly over the orbit, i.e. on a critical ratio:

$$\eta \equiv \frac{t_{\text{heat}}}{t_{\text{orb,c}}} \approx 1.6 \left(\frac{M}{5\, M_\odot}\right)^{-2/5} \left(\frac{Q_{\text{tot}}}{3\, \text{MeV}}\right)^{3/5} \left(\frac{t_{\text{heat}}}{1\, \text{s}}\right)^{2/5}. \quad (8)$$

If $t_{\text{orb}} \ll t_{\text{orb,c}}$ ($\eta \gg 1$), then the heating is applied in a short burst uniformly to all fluid elements and there is no runaway (preferential heating) of fluid elements as discussed above. In this case, no significant energy gap is opened; $M_{\text{fb}}$ shows a slight dip around the time at which $t_{\text{orb}} \sim t_{\text{heat}} \sim 1\, \text{s}$, but otherwise experiences no significant interruption of fall-back activity and $M$ still approaches a power-law $\propto r^{-5/3}$ decay at later times.

If $t_{\text{orb}} \gtrsim t_{\text{orb,c}}$ ($\eta \approx 1$) then an energy gap is opened in the debris. If $t_{\text{heat}} \gg t_{\text{orb,c}}$ ($\eta \gg 1$), then only the most marginally bound matter will receive enough heat to unbind before arriving at apocenter and the energy gap extends to $E > 0$. This case produces an absolute cut-off in $dM/dE$ (and hence $M_{\text{fb}}$) for $|E_{\text{tot}}| \lesssim E_{\text{c}}$, where $E_{\text{c}}$ is the energy of orbits of period $t_{\text{orb}} = t_{\text{orb,c}}$.

In intermediate cases for which $t_{\text{orb}} \sim t_{\text{orb,c}}$ ($\eta \sim 1$), a large gap is opened in the energy distribution of the debris, but now material with initial orbital periods $t_{\text{orb}} \lesssim t_{\text{orb,c}}$ remains marginally bound despite the extra energy it receives, thus opening up a large temporal gap in $M_{\text{fb}}$ (see Fig. 4 for an illustration). Specifically, we find that a critical value of $\eta = \eta_c \approx 1.25$ is needed to generate a long cut-off of $\sim 30\, \text{s}$, similar to the observed lulls in the SGRB EE light curves (Fig. 1). As we discuss later, the black hole mass dependence of $\eta_c$ may have implications for distinct fall-back behaviour in NS–BH versus NS–NS mergers.

Fig. 6 shows the results of a broader study of the outcomes of fall-back across the parameter space of $r$-process heating parameters $1$...
Figure 6. Regimes of the impact of $r$-process heating on fall-back accretion in the space of the total nuclear released energy, $Q_{\text{tot}}$, and the time-scale of the heating, $t_{\text{heat}}$. The value of $Q_{\text{tot}}$ (left vertical axis) is mapped on to the initial electron fraction $Y_e$ (right vertical axis) using equation (3) as shown in the bottom panel of Fig. 2. A 2D coloured histogram shows the distribution of particle heating trajectories ($Q_{\text{tot}}$, $t_{\text{heat}}$), from the SkyNet calculations (Fig. 2); since each particle has the same mass, this is also the mass distribution. The symbols show the results of our parameter study in which we assume a BH mass $M \approx 6M_\odot$ and that all fluid elements experienced heating characterized by fixed values of $Q_{\text{tot}}$ and $t_{\text{heat}}$ according to equation (4; see Fig. 5 for a few examples). The crosses denote cases that result in a complete cut-off in the fall-back rate after a given time, while the circles show cases in which a temporal gap is opened in the fall-back curve (the duration of the gap is denoted both by the radius and colour of the circle using the legend given on the right). The lines represent the critical condition ($\eta = \eta_c \approx 1.25$; equation 8) giving rise to a long ($\sim 30$ s) gap for different values of the mass of the central black hole as marked, ranging from $M \approx 2.5M_\odot$ relevant to NS–NS mergers to $M \approx 6–20 M_\odot$ relevant to NS–BH mergers. While the tidal ejecta from NS–NS mergers lies in the cut-off regime ($\eta \gg 1$) for neutron-rich ejecta $Y_e \lesssim 0.2$, the NS–BH merger case resides close to the gap regime ($\eta \sim \eta_c \approx 1.25$).

Our calculations shown in Fig. 5 were performed under the assumption that all fluid elements experienced heating characterized by single values of $Q_{\text{tot}}$ and $t_{\text{heat}}$. While reasonable as a first-order approximation (Fig. 2), in detail these parameters will vary between fluid elements as a result of a finite spread in their initial electron fraction $Y_e$ and precise thermodynamic conditions (e.g. entropy and expansion rate, which affect the properties of the seed nuclei). It is

$\leq Q_{\text{tot}} \leq 4$ MeV and $0 \leq t_{\text{heat}} \leq 3$ s. The crosses denote runs resulting in cut-off behaviour, while the circles denote cases with gaps (with the duration of gap indicated by the size of the circle and its colour, based on the key given on the right of the diagram). To highlight the relevant region of parameter space, we overlay with a coloured histogram the mass-weighted distribution ($Q_{\text{tot}}$, $t_{\text{heat}}$), obtained by mapping the ejecta properties from our NS–BH simulation data into the parameters extracted from the SkyNet heating trajectories (Lippuner et al. 2017) based on their $Y_e$ values (Fig. 2), where we define $t_{\text{heat}}$ as the time at which the heating rate curve first decreases below half of its maximum value. We also overlay with lines the condition $\eta = \eta_c$ (equation 8) for different assumptions about the BH masses, ranging from low values $M \approx 2.5M_\odot$ relevant to the remnants of NS–NS mergers to higher values $M \sim 6–20 M_\odot$ appropriate to NS–BH mergers. This location of the crosses and the circles relative to the $\eta = \eta_c$ line for the BH mass $M \approx 6M_\odot$ corresponding to our simulation verifies the validity of this criterion as that responsible for separating cut-off from gaps in fall-back behaviour.

3.1 $Y_e$-dependent spread in fluid element heating

Our calculations shown in Fig. 5 were performed under the assumption that all fluid elements experienced heating characterized by single values of $Q_{\text{tot}}$ and $t_{\text{heat}}$. While reasonable as a first-order approximation (Fig. 2), in detail these parameters will vary between fluid elements as a result of a finite spread in their initial electron fraction $Y_e$ and precise thermodynamic conditions (e.g. entropy and expansion rate, which affect the properties of the seed nuclei). It is
finite heating spread is to smooth out, but not eliminate, the gap in mass fall-back, in other words turning the ‘gap’ into a ‘lull’. Our initial conclusion that r-process heating can lead to at least partial gaps in the fall-back when the critical condition \( \eta \approx \eta_c \) is satisfied (where now \( \eta \) is defined using the mass-averaged values of \( Q_{\text{tot}} \) and \( t_{\text{heat}} \)) thus appears to be robust.

Our calculations thus far have employed a step-function heating profile as given by equation (4). To explore the sensitivity of our conclusions to this assumption, we perform an identical calculation where we directly use the direct heating curves from SkyNet (Fig. 2, top panel), which have been mapped on to the ejecta from our simulation data according to their closest \( Y_e \) value (e.g. 3). We still assume that heating for a given fluid element goes to zero when re-compressing (\( v_r < 0 \)). The results of this simulation, as shown in the bottom panel of Fig. 7, are fall-back with a gap of about 10 s. This is because the heating parameters for our fiducial simulation of an NS merging with a \( \approx 6 \, M_\odot \) BH overlaps with the gap condition \( \eta \approx \eta_c \) (the red solid line in Fig. 6).

### 4 IMPLICATIONS FOR EXTENDED EMISSION IN SGRBS

As described in Section 3, the condition \( \eta \approx \eta_c \approx 1.25 \) separates two distinct regions in the space of \( Q_{\text{tot}} - t_{\text{heat}} \) shown in Fig. 6: the lower left corner (\( \eta \ll \eta_c \)), where fall-back has a gap or is uninterrupted, and the upper right (\( \eta \gg \eta_c \)), where fall-back exhibits a complete cut-off.

The heating properties of the ejecta from our NS–BH simulation, which left a black hole of mass \( M \sim 6 \, M_\odot \), lie in the gap region close to the solid red line. We have confirmed this behaviour by calculating the fall-back rate directly using SkyNet heating trajectories (Fig. 7, bottom panel). However, because \( \eta \propto M^{-25} \) (equation 8), otherwise similar ejecta from a merger that resulted in central black hole of lower or greater mass would instead result in \( \eta \gg \eta_c \) or \( \eta \ll \eta_c \), respectively, and thus would exhibit qualitatively different fall-back behaviour.

Fig. 8 shows the condition \( \eta = \eta_c \), now in the space of black hole mass \( M \) and electron fraction \( Y_e \), where we have used the relationship \( Q_{\text{tot}}(Y_e) \) from equation (3; see bottom panel of Fig. 2).

If the X-ray luminosity of the extended prompt emission following a short GRB is proportional to the mass fall-back rate, \( L_X \propto \dot{M} \), then this reasoning would suggest that NS–NS mergers (\( M \approx 2.5 \, M_\odot \)) with similar ejecta properties would lie in the regime \( \eta \gg \eta_c \) and thus should generate little or no late-time fall-back and hence would not be accompanied by luminous extended X-ray emission. By contrast, NS–BH mergers, given their more massive BHs \( \gtrsim 5 \, M_\odot \), lie in the regime \( \eta \sim \eta_c \) and thus could produce fall-back with temporal gaps extending up to tens of seconds, in agreement with those short GRBs showing EE (Fig. 1).

One caveat is that, while these conclusions hold for highly neutron-rich ejecta (\( Y_e \lesssim 0.2 \)), as characterizes the equatorial tidal tails in NS–NS and NS–BH mergers, it would not necessarily apply to the most polar-concentrated shock-heated dynamical ejecta, which experiences weaker weak interactions. At least in the case of NS–NS mergers, this has been shown to give rise to a wider range of \( Y_e \) (Wanajo et al. 2014; Goriely et al. 2015; Sekiguchi et al. 2016), potentially extending to values \( Y_e \gtrsim 0.2 \) that would place even NS–NS mergers into the \( \eta \sim \eta_c \) regime and result in some fall-back. However, because in many cases the quantity of high \( Y_e \) matter is likely to be much less than the total, the amount of fall-back from this component would be significantly smaller and the presence of a cut-off in the fall-back rate might be preserved.
Our results suggest the possibility that the two apparent classes of SGRBs — those with and those without EE — may be associated with NS–BH and NS–NS mergers, respectively. Such a dichotomy of origin was previously suggested on the completely different basis of the observed distribution of spatial offsets of short GRBs from their host galaxies (Troja et al. 2008); however, the statistical significance of this difference was subsequently challenged (Fong & Berger 2013).

Is such a model consistent with current event rate constraints? A fraction $f_{\text{EE}} \gtrsim 0.2$–0.4 of SGRBs is accompanied by EE (Norris & Bonnell 2006). For our progenitor dichotomy scenario to hold, the ratio of the volumetric rate of NS–BH mergers to that NS–NS mergers must be at least as high as $f_{\text{EE}}/f_{\text{SGRB}}$, where we have assumed that all NS–NS mergers are accompanied by an SGRB but only a fraction $f_{\text{SGRB}} < 1$ of NS–BH mergers does the same (as the latter requires rapid BH spin in the prograde direction relative to the orbit for the NS to be tidally disrupted outside the BH horizon).

From the first and second observing runs of Advanced LIGO, the observed rate of NS–NS mergers is $R_{\text{NSNS}} = 110 - 3840 \text{Gpc}^{-3} \text{yr}^{-1}$ at 90 per cent confidence, while the upper limit on the rate of mergers of $\gtrsim 5 \text{M}_\odot$ with NSs is $R_{\text{NSBH}} \gtrsim 610 \text{Gpc}^{-3} \text{yr}^{-1}$ (The LIGO Scientific Collaboration & the Virgo Collaboration 2018). The ratio of these rates is thus only weakly constrained at present, such that as long as a moderate fraction of NS–BH mergers produce GRBs,

$$f_{\text{SGRB}} \gtrsim f_{\text{EE}} \left( \frac{R_{\text{NSNS}}}{R_{\text{NSBH}}} \right) \gtrsim 0.05 \left( \frac{f_{\text{EE}}}{0.3} \right),$$

one cannot yet rule out the possibility that NS–BH mergers are sufficiently common to account for the population of SGRBs with EE ($f_{\text{EE}} \gtrsim 0.2$–0.4). While thus far one NS–NS merger has been discovered compared to zero NS–BH mergers, the statistics of the current sample are obviously very small. If all EE is attributed to fall-back accretion in NS–BH mergers, then our model implies that, depending on the future sensitivity of gravitational wave detectors, the steady-state discovery rate of NS–BH mergers could be significantly higher than that of NS–NS mergers (and that a sizable fraction of the BHs in these systems is spinning in the prograde orbital direction).

5 CONCLUSIONS

Despite the recent discovery of a short burst of gamma-rays in association with the gravitational waves from an NS–NS merger, the origin of the temporally extended X-ray emission that is observed following a significant fraction of short GRBs remains a mystery. Late-time activity from the black hole accretion disc powered by fall-back of marginally bound debris, as would naïvely be expected in both NS–NS and NS–BH mergers, has been proposed as a source of this behaviour (e.g. Rosswog 2007). Following Metzger et al. (2010a), we have employed a simple model to explore the impact of $\beta$-decay heating due to $r$-process nucleosynthesis on the time dependence of the mass fall-back rate, using initial data on the ejecta fluid elements (Fig. 3), and the properties of the $r$-process heating received (Fig. 2), extracted directly from NS–BH merger simulations.

We confirm that this $r$-process heating significantly alters the fall-back rate from the canonical $M \propto t^{-5/3}$ behaviour, generating instead either an abrupt cut-off, or temporal gap, in the fall-back rate on a time-scale of $\sim 10-100$ s (Figs 5 and 6). This behaviour is robust to the presence of a realistic spread in the heating properties of the fluid elements imparted by a realistic range in the electron fraction and thermodynamic history (Fig. 7). Whether a cut-off or gap behaviour is obtained depends on the value of a critical dimensionless parameter $\eta$ (equation 8). The dependence of $\eta$ on black hole mass suggests a possible distinction between cut-off behaviour in NS–NS mergers (low black hole mass) and delayed fall-back in NS–BH mergers (high black hole mass), as illustrated in Fig. 8. The presence or absence, respectively, of EE thus provides a possible way to distinguish NS–BH from NS–NS mergers.

Our model could be improved or extended along several fronts in future work. In addition to the BH mass, the critical parameter $\eta$ depends on the total energy $Q_{\text{tot}}$ and heating time-scale $t_{\text{heat}}$ of the $r$-process. These properties are related to the $Q$ values and $\beta$-decay rates of neutron-rich isotopes whose masses and other properties have yet to be measured by laboratory experiments (e.g. Horowitz et al. 2018 and references therein). A more thorough parameter study of the range of $t_{\text{heat}}$ and $Q_{\text{tot}}$, e.g. assuming different theoretical models for the nuclear masses, would provide an additional check on the robustness of our conclusions.

More ambitiously, multidimensional hydrodynamical simulations of the fall-back process, accounting for the effects of $r$-process heating, are needed for a more robust assessment of the gap or cut-off formation process. Our treatment of directly placing the $r$-process thermal energy into debris kinetic energy neglects the transfer of thermal energy between adjacent fluid elements, though the highly supersonic and nearly radial motion of the debris should mitigate these effects (Fig. 4). Models that do not self-consistently include the backreaction of the thermodynamics of the fluid elements on the $r$-process path (Rosswog et al. 2014) may be inadequate because once matter starts to recompress and adiabatically heat, the rate of $\beta$-decay heating will be substantially suppressed as the $r$-process path moves back towards the stable valley due to the higher temperatures (Metzger et al. 2010a). This complex 'feed-back' process may ultimately necessitate coupling
at least a simplified r-process network (e.g. a one-zone model such as that of Lattimer et al. 1977) directly into the hydrodynamical simulations.

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