On the astrophysical robustness of the neutron star merger r-process

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ABSTRACT
In this study we explore nucleosynthesis in the dynamic ejecta of compact binary mergers. We are particularly interested in the question how sensitive the resulting abundance patterns are to the parameters of the merging system. Therefore, we systematically investigate combinations of neutron star masses in the range from 1.0 to 2.0 M⊙ and, for completeness, we compare the results with those from two simulations of a neutron star black hole merger. The ejecta masses vary by a factor of 5 for the studied systems, but all amounts are (within the uncertainties of the merger rates) compatible with being a major source of the cosmic r-process. The ejecta undergo robust r-process nucleosynthesis which produces all the elements from the second to the third peak in close-to-solar ratios. Most strikingly, this r-process is extremely robust, and all 23 investigated binary systems yield practically identical abundance patterns. This is mainly the result of the ejecta being extremely neutron rich (Y_e ≈0.04) and the r-process path meandering along the neutron drip line so that the abundances are determined entirely by nuclear rather than astrophysical properties. While further questions related to galactic chemical evolution need to be explored in future studies, we consider this robustness together with the ease with which both the second and third peak are reproduced as strong indications that compact binary mergers are prime candidates for the sources of the observed unique heavy r-process component.

Key words: equation of state – gravitation – hydrodynamics – neutrinos – nuclear reactions, nucleosynthesis, abundances.

1 INTRODUCTION
About half of the elements heavier than iron are formed by neutron capture reactions that occur rapidly in comparison with β-decays. The basic physical mechanisms of this ‘rapid neutron capture’ or ‘r-process’, for short, had already been identified in the seminal paper of Burbidge et al. (1957). Nevertheless, the cosmic coudlrons in which these heavy elements are forged have remained elusive for more than half a century. Observations of metal-poor stars point to at least two groups of r-process events. The first one is rare and produces predominantly lighter elements from strontium to silver (Cowan & Sneden 2006; Honda et al. 2006). Given that its signature is less unique the second component might also be the result of a superposition from several sources (see e.g. Arcones & Montes 2011; Frischknecht, Hirschi & Thielemann 2012). Metal-poor stars that are enriched by r-process material show patterns that very closely match the (scaled) Solar system abundances for nuclei beyond Ba (Z = 56). This suggests that this heavy r-process component is produced only if a very unique set of astrophysical conditions is realized, or, alternatively, that it is produced in a range of realizations, but is insensitive to the exact parameters of its formation environment. So far, there is no generally accepted explanation for the unique heavy r-process component.

Historically, supernovae were considered the ‘natural’ source of r-process elements, but this view has recently been distressed by a slew of investigations (e.g. Arcones, Janka & Scheck 2007; Fischer et al. 2010; Hudepohl et al. 2010; Roberts, Woosley & Hoffman 2010). These studies found neutrino-driven winds in supernovae to be seriously challenged in providing the physical conditions (high entropy, low electron fraction together with rapid expansion) that are required to produce the heavy (A > 90) r-process elements. A possible exception may be magnetorotationaly driven supernova (SN) jets where the fast ejection of highly compressed neutron-rich

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matter produces interestingly low electron fraction values (Winteler et al. 2012). It remains to be explored, however, how robust this scenario is with respect to the stellar parameters and with respect to its nucleosynthetic yields. The main contenders of supernovae in terms of r-process nucleosynthesis are compact binary mergers of either two neutron stars (ns2) or a neutron star and a black hole (nsbh; Lattimer & Schramm 1974, 1976; Eichler et al. 1989; Freiburghaus, Rosswog & Thielemann 1999).

Since the first nucleosynthesis calculations for ns merger ejecta (Freiburghaus et al. 1999) were performed, numerical models have seen vast improvements in physics, numerical methods and resolution, and we consider it timely to address this important topic again. There have been recent r-process studies in a compact binary merger context (Goriely, Bauswein & Janka 2011; Roberts et al. 2011; Wanajo & Janka 2012) which, in agreement with the results of the first study, confirm the occurrence of a robust r-process. While all previous studies provided new insights each of them had its own limitations. For example, in all of them the ejecta electron fraction was poorly known since Y_e-changing weak interactions were ignored. More importantly, the ns2 parameter space was explored only very punctually around ns masses of 1.4 M⊙ (Freiburghaus et al. 1999) explored one case from Rosswog et al. (1999), two cases were explored in the conformal flatness approximation approach of Goriely et al. (2011) and three cases in Roberts et al. (2011). In recent years, however, it has become clear that the neutron stars realized in nature span a substantial range in masses, see e.g. Lattimer & Prakash (2010) for compilation of ns masses. Therefore, an important open question needs to be answered: ‘How sensitive is the compact object merger nucleosynthesis to the actual astrophysical system that merges?’

This is what we address in this paper: we systematically scan the ns binary parameter space in 21 simulations and calculate the resulting nucleosynthetic products. For completeness, we also show the results for the ejecta of two nsbh systems. Apart from exploring a very large ns mass range, we also significantly improve on the treatment of the electron fraction Y_e: we start with cold neutron stars in β-equilibrium and allow for the evolution of Y_e via electron and positron capture (Rosswog & Liebendörfer 2003). We demonstrate that the nucleosynthesis results are essentially independent of the parameters of the merging system, i.e. every ns2 or nsbh system produces nearly exactly the same abundance pattern. The results show, however, some sensitivity to the properties of extremely neutron-rich nuclei (e.g. binding energy, half-lives, fission properties) which are not well known (Thielemann et al. 2011).

Although heavy element nucleosynthesis is the clear focus of this study, we would like to stress the importance of compact object merger ejecta and their nucleosynthesis in a broader astrophysical perspective. Both the LIGO and Virgo gravitational wave (GW) detectors are currently being upgraded (Abbott et al. 2009; Smith 2009; Sengupta et al. 2010) to an ~10–15 times better sensitivity than the original versions of the instruments. This will increase the volume of accessible astrophysical sources by more than a factor of 1000 reaching a detection horizon of a few hundred Mpc for ns2 mergers and about a Gpc for nsbh mergers (Abadie et al. 2010). Since the first detections will most likely be around or even below threshold, it is of paramount importance to identify accompanying electromagnetic signatures of the most promising GW sources to enhance the confidence in a possible detection. Short gamma-ray bursts are likely the brightest electromagnetic manifestations that result from a compact binary merger, but also the r-process nucleosynthesis, and the decay of the resulting radioactive nuclei may leave an observable signature as optical/ultraviolet transients, so-called ‘macronovae’ (Li & Paczynski 1998; Kulkarni 2005; Rosswog 2005; Metzger et al. 2010b; Roberts et al. 2011; Metzger & Berger 2012). For ns mergers they are expected to peak about 10 h after the mergers with a few times 10^{42} erg s⁻¹ (Piran, Nakar & Rosswog 2012; Rosswog, Piran & Nakar 2012). Later, on timescales of months to years, the ejecta dissipate their kinetic energy in the ambient medium and thereby produce possibly detectable radio flares (Nakar & Piran 2011; Piran et al. 2012).

This paper is structured as follows. In Section 2 we briefly summarize our simulations and explore the amount of ejected matter as a function of the binary system parameters. In Section 3 we describe how we calculate the nuclear abundances and discuss how they depend on the astrophysical properties of the merging system and on the nuclear physics input. In Section 4 we summarize our results and discuss their astrophysical implications.

2 HYDRODYNAMIC EVOLUTION AND DYNAMICALLY EJECTED MASS

2.1 Explored parameter space

Until recently, the ns mass distribution was thought to be clustered narrowly around 1.35 M⊙ (Thorsett & Chakrabarti 1999). Therefore, essentially all ns merger studies have focused on a narrow range of masses around this value. Recent observations, however, indicate a much broader ns mass spectrum. There is now ample support for ns masses significantly larger than 1.5 M⊙. A broad peak around 1.5–1.7 M⊙ has been found (Kiziltan, Kottas & Thorsett 2010; Valentim, Rangel & Horvath 2011) for neutron stars with white dwarf companions, and an additional low-mass peak near 1.25 M⊙ is thought to be characteristic for neutron stars that were produced by electron capture supernovae (Podsiadlowski et al. 2004; van den Heuvel 2004; Schwab, Podsiadlowski & Rappaport 2010). PSR J1614−2230 with 1.97 ± 0.04 M⊙ possesses the largest accurately known ns mass (Demorest et al. 2010), but even higher masses may exist in nature. For example, the mass of the black widow pulsar has recently been determined as 2.4 ± 0.12 M⊙ (van Kerkwijk, Breton & Kulkarni 2011). On the lower mass side, the secondary ns in J1518+4904 has a best estimated mass value of only 0.72 M⊙, although with a very large 1σ error bar of 0.5 M⊙, see Lattimer & Prakash (2010) and references therein.

We take these results as a motivation for a systematic exploration of the parameter space from 1.0 to 2.0 M⊙ in steps of 0.2 M⊙. Since the ns viscosity cannot substantially spin up the neutron stars during the short tidal interaction phase preceding the merger (Bildsten & Cutler 1992; Kochanek 1992), all our models have vanishing initial ns spins. Our ns2 cases are complemented by two nsbh cases with black hole masses of 5 and 10 M⊙; for an overview over the performed simulations see Table 1.

2.2 Methodology

To follow the hydrodynamic evolution we use the smooth particle hydrodynamics (SPH) method (see Monaghan 2005 and Rosswog 2009 for recent reviews), which, due to its completely Lagrangian nature, is ideally suited to follow the ejected material. Our code is an updated version of the one that was used in earlier studies

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1 These authors also used the Lattimer–Swesty equation of state (EOS) to test for the impact of the ns EOS.
only ejects small amounts of the i-component, while the fraction as 'tidal' or t-component. As the mass ratio decreases, the primary loses more mass to the t-component. Oechslin, Janka & Marek (2007) found that the ejecta stem from two distinct regions: a hot component from the interaction region between the stars (subsequently called 'interaction-' or i-component for short) and a colder one where matter is simply flung out by tidal torques, see the first two panels in Fig. 2. We refer to this matter as 'tidal' or t-component. As the mass ratio decreases, the primary only ejects small amounts of the i-component, while the fraction and overall amount in the t-component from the secondary increase. This differs from the findings of Goriely et al. (2011), where even in unequal-mass cases most of the mass is lost from the primary as the t-component. For now we can only speculate about the origin of this difference: it could either be due to the different treatment of self-gravity (conformal flatness versus Newtonian) or, since the i-component comes from a region of very strong shear, it could also to some extent be impacted by the specifics of the numerical method, such as implementation of artificial viscosity.

As can be seen from the volume rendering in panel 3 of Fig. 2, the ejecta are extremely neutron rich. In Fig. 3, left-hand panel, we bin the mass fractions of the ejecta of all simulations according to the electron fraction. The values are clustered around Y_e ∼ 0.04 ± 0.02, ρ = (1.4 ± 0.5) × 10^{14} g cm^{-3}, corresponding to the outer ∼ 2 km of the original ns (see, e.g., fig. 1 in Rosswog et al. 2012), but also continues to expand adiabatically and without shocks. Therefore, all fluid elements in tidal ejecta share approximately the same thermodynamic history and provide robust conditions for uniform nucleosynthesis.

2.3 Ejecta properties

The merger dynamics is rather sensitive to asymmetries in the component masses, see Fig. 1 (for a more complete overview over the remnant structures see fig. 1 in Piran et al. 2012). The asymmetries enhance the degree to which the secondary is disrupted and they lead to larger ejecta masses and velocities. As already pointed out by Oechslin, Janka & Marek (2007), the ejecta stem from two different regions: a hot component from the interaction region between the stars (subsequently called 'interaction-' or i-component for short) and a colder one where matter is simply flung out by tidal torques, see the first two panels in Fig. 2. We refer to this matter as 'tidal' or t-component. As the mass ratio decreases, the primary only ejects small amounts of the i-component, while the fraction

| Run | m_1 (M_⊙) | m_2 (M_⊙) | N_{SPH} (10^6) | t_{end} (ms) | m_{ej} (M_⊙) | Comment |
|-----|-----------|-----------|----------------|-------------|--------------|---------|
| 1   | 1.0       | 1.0       | 1.0            | 15.3        | 7.64 × 10^{-3} | Secondary still orbiting |
| 2   | 1.2       | 1.0       | 1.0            | 15.3        | 2.50 × 10^{-2} | Secondary still orbiting |
| 3   | 1.4       | 1.0       | 1.0            | 16.5        | 2.91 × 10^{-2} | Secondary still orbiting |
| 4   | 1.6       | 1.0       | 1.0            | 31.3        | 3.06 × 10^{-2} | Secondary still orbiting |
| 5   | 1.8       | 1.0       | 1.0            | 30.4        | >1.64 × 10^{-2} | Secondary still orbiting |
| 6   | 2.0       | 1.0       | 0.6            | 18.8        | >2.39 × 10^{-2} | Secondary still orbiting |
| 7   | 1.2       | 1.2       | 1.0            | 15.4        | 1.68 × 10^{-2} | Secondary still orbiting |
| 8   | 1.4       | 1.2       | 1.0            | 13.9        | 2.12 × 10^{-2} | Secondary still orbiting |
| 9   | 1.6       | 1.2       | 1.0            | 14.8        | 3.33 × 10^{-2} | Secondary still orbiting |
| 10  | 1.8       | 1.2       | 1.0            | 21.4        | 3.44 × 10^{-2} | Secondary still orbiting |
| 11  | 2.0       | 1.2       | 0.6            | 15.1        | >2.95 × 10^{-2} | Secondary still orbiting |
| 12  | 1.4       | 1.4       | 1.0            | 13.4        | 1.28 × 10^{-2} | Secondary still orbiting |
| 13  | 1.6       | 1.4       | 1.0            | 12.2        | 2.36 × 10^{-2} | Secondary still orbiting |
| 14  | 1.8       | 1.4       | 1.0            | 13.1        | 3.84 × 10^{-2} | Secondary still orbiting |
| 15  | 2.0       | 1.4       | 0.6            | 15.0        | 3.89 × 10^{-2} | Secondary still orbiting |
| 16  | 1.6       | 1.6       | 1.0            | 13.2        | 1.97 × 10^{-2} | Secondary still orbiting |
| 17  | 1.8       | 1.6       | 1.0            | 13.0        | 1.67 × 10^{-2} | Secondary still orbiting |
| 18  | 2.0       | 1.6       | 0.6            | 12.4        | 3.79 × 10^{-2} | Secondary still orbiting |
| 19  | 1.8       | 1.8       | 1.0            | 14.0        | 1.50 × 10^{-2} | Secondary still orbiting |
| 20  | 2.0       | 1.8       | 0.6            | 11.0        | 1.99 × 10^{-2} | Secondary still orbiting |
| 21  | 2.0       | 2.0       | 0.2            | 21.4        | 1.15 × 10^{-2} | Secondary still orbiting |
| 22  | 5.0×      | 1.4       | 0.2            | 138.7       | 2.38 × 10^{-2} | nsbh     |
| 23  | 10.0×     | 1.4       | 0.2            | 139.0       | 4.93 × 10^{-2} | nsbh     |

(Rosswog & Davies 2002; Rosswog & Liebendörfer 2003; Rosswog, Ramirez-Ruiz & Davies 2003; Rosswog 2005). It uses the Shen et al. EOS (Shen et al. 1998a,b) and an opacity-dependent multi-flavour neutrino leakage scheme (Rosswog & Liebendörfer 2003). The latter allows us in particular to follow the Y_e changes due to electron and positron captures. Moreover, we use a modern, time-dependent artificial viscosity prescription; see Rosswog et al. (2000, 2008) for details. The presented calculations make use of Newtonian gravity which leads to less compact neutron stars than General Relativity (GR). This could have a moderate effect on the amount of ejecta mass and, since the β-equilibrium Y_e is determined by the density inside the initial ns, the electron fraction. As we will show below, however, the heavy nucleosynthesis is only affected once Y_e reaches ∼ 0.2 and the ejecta masses are (within the existing uncertainties) still consistent with compact binary mergers being a major r-process source even if the amount of ejecta is reduced by a factor of a few. Therefore, we expect our major results to be robust. Nevertheless this point needs improvement in future work.
Figure 1. Sensitivity to mass ratio: shown are density cuts ([600 km × 600 km]), colour-coded is the logarithm of density in cgs units] of a 1.2–1.4 M⊙ (t = 13.9 ms), a 1.4–1.4 M⊙ (t = 13.4 ms) and a 2.0–1.4 M⊙ (t = 15.0 ms) merger.

Figure 2. The ejecta come from two different regions (left and middle): a hot interaction region between the stars where matter is ejected by hydrodynamic effects and a colder region that is flung out by tidal torques. The fraction of the latter material increases with the asymmetry in the stellar masses (left: 1.4 + 1.4 M⊙, right: 1.8 + 1.4 M⊙). Right: volume rendering of the Y_e distribution (1.4 + 1.2 M⊙ at t = 8.09 ms); only matter below the orbital plane is shown.

Figure 3. Left-hand panel: distribution of Y_e in the selected subset of ejected particles in all ns^2 and nsbh simulations, binned by mass. The inset shows how Y_e is correlated with the initial density in each particle before the merger. Majority of points on the Y_e–ρ diagram trace the ns composition in the layer below the crust, where the Y_e(ρ) reaches minimum. As a consequence, Y_e in the ejecta is narrowly distributed around Y_e ≈ 0.04. Right-hand panel: fit of the ejected mass normalized by the total mass of the binary as a function of the asymmetry parameter η (see the main text for definition). Coefficients of the fit are A = 0.0125, B = 0.015, C = 0.0083 and σ = 0.0056. The arrows indicate the lower limits on the ejected mass for the three simulations in which the secondary is still not fully disrupted.
systems no unique trend is found with respect to the total binary mass.

For at least small ns mass asymmetries, i.e. for the most common cases, we provide a fit formula based on the dimensionless mass asymmetry parameter
\[ \eta := 1 - 4m_1m_2/(m_1 + m_2)^2. \] (1)
Contrary to the mass ratio \( q \), \( \eta \) is symmetric with respect to both masses and varies only in a finite range from 0 to 1 (with 0 for an equal-mass system and \( \eta \) approaching 1 for extreme mass ratios).²
The fit formula
\[ m_\omega(m_1, m_2) = (m_1 + m_2) \left( A - B\eta - C / (1 + \eta^2/\sigma^2) \right), \] (2)
with \( A = 0.0125 \), \( B = 0.015 \), \( C = 0.0083 \), and \( \sigma = 0.0056 \), as shown as a solid line in Fig. 3, right-hand panel, provides a reasonably good approximation to the simulation results. For unequal-mass binaries \( m_\omega/(m_1 + m_2) \) tends to a minimum when approaching \( \eta \to 0 \) and a maximum of \( \sim 0.12 \) is reached near \( \eta = 0.02 \).

3 NUCLEOSYNTHESIS
3.1 Hydrodynamic trajectories
Our nucleosynthesis calculations are performed with a large reaction network (Winteler 2012; Winteler et al. 2012) that is based on the BasNet network (Thielemann et al. 2011). It includes over 5800 isotopes from nucleons up to \( Z = 111 \) between the neutron drip line and stability. The reaction rates are from the compilation of Rauscher & Thielemann (2000) for the finite range droplet model (FRDM; Möller et al. 1995) and the weak interaction rates (electron/positron captures and \( \beta \)-decays) are the same as in Arcones & Martinez-Pinedo (2011). In addition, neutron capture and neutron-induced fission rates of Panov et al. (2010) and \( \beta \)-delayed fission probabilities as described in Panov et al. (2005) are used.

The thermodynamical conditions are taken from hydrodynamical trajectories of individual SPH particles. We calculate nucleosynthetic yields for 30 representative trajectories in each of the 21 ns² merger simulations, and for 20 trajectories in the two nsbh merger simulations. These trajectories are a fair representation of the total ejected mass, see Table 1.

Once matter has been ejected from the central high-density region, both temperature and density decrease rapidly while the electron fraction \( Y_e \) can, in principle, start to increase due to \( \beta \)-decays and positron captures. A small fraction of the trajectories (less than 0.008, see Fig. 3, left-hand panel) suffers shocks during their later evolution during which their \( Y_e \) increases up to \( \sim 0.25 \), but for the bulk of trajectories the \( Y_e \)-values remain essentially constant. Some of the higher \( Y_e \)-material also comes from the (marginally resolved) ns crust material. Although heavily dominated by values around 0.03, our ejecta electron fractions show a broader range than those of Goriely et al. (2011), probably due to our higher resolution (crust material) and their omission of weak interactions.

We start our nucleosynthesis calculations when the temperature of the trajectory has dropped below \( T_0 = 10 \) GK. The initial composition is given by nuclear statistical equilibrium and it is further evolved with the complete network. Because most of the trajectories from hydrodynamical simulations terminate after \( t_{\text{fin}} \sim 10-20 \) ms, we extrapolate according to a free expansion for the density and an adiabatic expansion law for the temperature:
\[ \rho(t) = \rho_{\text{fin}} \left( \frac{t}{t_{\text{fin}}} \right)^{-3}, \quad T(t) = T \left[ S_{\text{fin}}, \rho(t), Y_e(t) \right], \]
where \( \rho_{\text{fin}} \) and \( S_{\text{fin}} \) are the density and entropy at \( t_{\text{fin}} \), while the temperature is calculated at each point of time from the EOS (Timmes & Swesty 2000).

Contrary to neutrino-driven winds where the ejecta are dominated by \( \sigma \) particles and only few seed nuclei participate in the \( r \)-process, essentially all matter in ns mergers undergoes the \( r \)-process. Therefore, the \( r \)-process energy generation is non-negligible and leads to a substantial temperature increase (see, e.g., Freiburghaus et al. 1999; Metzger et al. 2010b; Goriely et al. 2011) that could affect the late-time evolution (Metzger et al. 2010b). We calculate the \( r \)-process heating in a post-processing step and consider its influence on the nucleosynthesis by modifying the temperature. To account for neutrino energy losses associated with \( \beta \)-decays, we introduce a heating efficiency parameter \( \epsilon_{\text{th}} \) which measures the fraction of nuclear power which is retained in the matter. Metzger et al. (2010a) argue that this fraction must be \( \epsilon_{\text{th}} \approx 0.25 - 1 \). As a default, we use \( \epsilon_{\text{th}} = 0.5 \), but below we explore how (in)sensitive the results are to this choice. For a given trajectory we take its initial entropy \( S_{\text{ini}}(t) \) and increment it by \( \epsilon_{\text{th}} S_{\text{fin}}(t) \), following the prescription of Freiburghaus et al. (1999).

3.2 Final abundances
The \( r \)-process nucleosynthesis in compact binary mergers is extremely robust; all 23 cases of Table 1 deliver essentially identical abundance distributions. We illustrate the origin of this robustness using the reference case of \( m_\odot \) with \( 1.4 M_\odot \) each (run 12 in Table 1). Fig. 4 shows the density and temperature evolution for all trajectories (upper panel) and their final abundances. Between different trajectories of a single run there is only a very small spread in the resulting abundances, although the evolution of temperature is not always the same. The resulting final abundances (sum over all ejecta trajectories for each run) are practically the same for all cases. We show a representative selection of the results from different runs in the bottom panel of Fig. 4.

The robustness of the abundances is mainly due to two factors: (1) the \( r \)-process path always reaches the drip line because of the extremely neutron-rich conditions and (2) there are several fission cycles (similar to what was observed in Goriely et al. 2011). These conditions lead to the evolution of the neutron density shown in Fig. 5, upper panel. The huge values of the neutron density are due to the (narrow range of) very low \( Y_e \)-values.

In our calculations the low \( Y_e \) results in an initial composition consisting of neutron-rich nuclei (\( Z > 20 \), on the neutron drip line) and neutrons. Differences with Goriely et al. (2011) in lighter heavy elements are not particularly relevant because these can also be produced in other astrophysical events, and the overall yields of the elements with \( A < 120 \) that we obtain are much smaller compared to the heavy robust \( r \)-process elements (similar to Goriely et al. 2011). Note that the contribution of the high-\( Y_e \) trajectories to the final abundances is negligibly small.

In order to understand the evolution of the \( r \)-process path we monitor the average neutron separation energy (middle panel of Fig. 5) defined as
\[ \langle S_{n} \rangle = \frac{\sum_{Z,A} S_n(Z,A)Y(Z,A)}{\sum_{Z,A} Y(Z,A)} \] (3)
Robust r-process in neutron star mergers

Figure 4. Top: density and temperature evolution for a bundle of trajectories, colour-coded by density at $T = 10$ GK. Middle: resulting final abundances distribution. Their averaged distribution is shown as a black dashed line, and a bold red line represents abundances for a trajectory without heating. All trajectories represent a subset from the standard 1.4–1.4 M⊙ merger. Bottom: distribution of abundances for a variety of different (ns and nsbh) merger cases. All different astrophysical systems yield essentially identical resulting abundances.

Figure 5. Evolution of the neutron density $N_n$ (top), average neutron separation energy $\langle S_n \rangle$ (middle) and average proton number $\langle Z \rangle$ (bottom) for five trajectories from run 12 with different initial densities (in the range $10^{12}–10^{13}$ g cm$^{-3}$).

with $S_n(Z, A)$ and $Y(Z, A)$ being the neutron separation energy and abundance of the nucleus $(Z, A)$. The average separation energy decreases when matter moves away from stability and it is, by definition, zero at the neutron drip line. The second panel of Fig. 5 shows that the average separation energy is initially below $\approx 1$ MeV, which indicates that the r-process path proceeds along the neutron drip line. The average proton number increases to $Z = 40$ at $t \approx 10^{-2}$ s where the neutron separation energy reaches a maximum. This local maximum occurs when the magic number $N = 82$ is overcome. Whenever the r-process path reaches a neutron magic number, it moves closer to the line of $\beta$-stability (i.e. larger $S_n$ values) by increasing $Z$ without changing $N$. After the matter flow passes $N = 82$ (here around $Z = 40$), the $\langle S_n \rangle$ decreases because the path again
gets further away from the $\beta$-stable region. This continues until the next magic number, $N = 126$, is encountered (corresponding to the minimum of $\langle S_n \rangle$ between 0.01 and 0.1 s). Around $t = 0.1s$ and $Z = 60$, $N = 126$ is also overcome as indicated by the second maximum of $\langle S_n \rangle$. After this point the oscillations in $\langle S_n \rangle$ are due mainly to fission as can be seen by the behaviour of $\langle Z \rangle$. Note that FRDM predicts a magic number at $N = 184$; consequently, the path reaching this point may lead to oscillations in $\langle S_n \rangle$. Therefore, the quantity $\langle Z \rangle$ is a nuclear fission indicator that is better suited than the average mass number (Goriely et al. 2011), since it can decrease only through fission reactions, while the average mass number can also decrease due to photodissociations. The maximum $\langle Z \rangle$ corresponds to the moment when a significant amount of matter has reached the region where fission becomes important. The daughter nuclei resulting from fission capture neutrons move first towards the drip line and then to higher $Z$ where fission acts again. In that way several fission cycles occur and lead to oscillations in $\langle Z \rangle$. In Fig. 5 one can distinguish at least three fission cycles. The final increase of $\langle S_n \rangle$ is due to the r-process freeze-out at $t \approx 1s$ when the neutron-to-seed ratio drops below unity and the matter $\beta$-decays to stability. We have used our reference case of two $1.4 M_\odot$ neutron stars, run 12, for illustration purposes, but behaviour is very similar in all of the other cases.

The major result of this study is that the r-process abundances for compact binary mergers do not depend on the astrophysical parameters of the merging system. They do depend, however, on the not-so-well-known properties of nuclei close to the neutron drip line. The most critical inputs are the nuclear masses, which determine the location of the drip line and the evolution of the path, and the fission barriers and yields distribution, which are responsible for the abundances within $120 < A < 170$. We explicitly explore how sensitive our results are to the details of our nuclear physics input. In Fig. 6, left-hand panel, we show the effect of including heating with the generated nuclear heat comes mainly from fission as can be seen by the behaviour of $\langle Z \rangle$. Note that the used prescription of Panov uses only two daughter nuclei while he also provides yield distributions in a more recent paper (Panov, Korneev & Thielemann 2008). Therefore, the deviation of our results from the Solar system abundances is sensitive to the physics of extremely neutron-rich nuclei.

In Fig. 7, left-hand panel, we show the evolution of a trajectory in the pressure–temperature plane for different heating efficiencies, the right-hand panel shows the generated nuclear power. The early evolution is very similar for all four cases because the internal energy is high and the contribution from the generated heat is negligible. The latter becomes the dominant part of the internal energy after about 20 ms when the temperature and density decrease enough for the r-process to begin. Similar to Freiburghaus et al. (1999) we find that the generated nuclear energy raises the temperature by almost an order of magnitude compared to the case without heating. Matter essentially stays in $(n, \gamma)$–$(\gamma, n)$ equilibrium until the neutron-to-seed ratio drops below unity ('freeze-out'). Note that all non-zero heating efficiencies produce trajectories that lie on the $P = 1/3aT^4$ line at freeze-out and therefore follow a similar behaviour in their decay to stability. The reason is that for our extremely neutron-rich conditions the main contribution to the pressure comes from neutrons and photons, while the electron pressure is orders of magnitude smaller. Therefore, at the freeze-out point when most of the neutrons are captured the system is radiation pressure dominated.

Later the system evolution enters the radiation-dominated expansion which was studied in Li & Paczyński (1998). In this regime the generated nuclear heat comes mainly from $\beta$-decays and can overall be well approximated by a power law, while before the freeze-out the generated heat stays approximately constant (see Fig. 7). We suggest the following fit which smoothly interpolates between a constant value and a power law:

$$\epsilon(t) = \epsilon_0 \left( \frac{1}{2} - \frac{1}{\pi} \arctan \frac{t - t_0}{\sigma} \right)^\alpha \times \left( \frac{\epsilon_0}{0.5} \right),$$

where $\epsilon(t)$ is the heating efficiency, $\epsilon_0$ is a threshold value, $t_0$ and $\sigma$ are fit parameters, and $\alpha$ is the power-law exponent. This fit smoothly transitions from a constant value at $t = t_0$ to a power law with exponent $\alpha$ at $t = t_0 + \sigma$. The choice of the parameters depends on the specific astrophysical scenario and can be determined through a detailed analysis of the system evolution.

At higher temperatures the radiation pressure must necessarily include the contribution from electron–positron pairs (Witti, Janka & Takahashi 1994; Farouqi et al. 2010).
The deviations from a power-law fit $\alpha t^{-\alpha}$ with index $\alpha = 1.3$. The small black triangles on each trajectory are neutron freeze-out points (the points at which $Y_e/Y_{\text{seed}} = 1$).

The $Y_e$ increase is triggered by shocks at low density, the neutron density drops earlier and the neutron captures freeze out.

**4 SUMMARY AND DISCUSSION**

In this study we have re-examined the question of heavy element nucleosynthesis in compact binary mergers. Our study adds several new aspects in comparison with earlier work. First, we systematically cover the plausible ns$^2$ parameter space with masses from 1.0 to 2.0 $M_\odot$ in 21 simulations. Despite the long history of ns$^2$ simulations we are not aware of any study that has systematically explored such a wide range of ns masses. We complement these cases with mergers of a 1.4 $M_\odot$ ns with black holes of 5 and 10 $M_\odot$. Secondly, since we use a nuclear EOS with neutron stars in initial cold $\beta$-equilibrium and include $Y_e$-changing weak interactions, we overcome the $Y_e$-ambiguity that has plagued all previous studies. We subsequently find the final abundances by following the nucleosynthesis along a large set of hydrodynamic trajectories with a state-of-the-art nuclear reaction network.

Consistent with earlier studies, we find a very robust r-process and final abundances that are in good agreement with the heavy Solar system abundance pattern. The major new result is that the final pattern is extremely robust across the whole parameter space: all 21 ns$^2$-merger and the two nsbh-merger cases yield practically identical nucleosynthesis outcomes. The major reason for this unique abundance pattern is the extreme neutron richness of the ejecta, $\langle Y_e \rangle \approx 0.04$. Consequently, in each case the r-process path meanders along the neutron drip line and matter undergoes several fission cycles, so that the abundances are determined entirely by nuclear rather than by astrophysical properties. As a corollary, the poorly known nuclear properties near the neutron drip line do have an impact on the resulting abundance pattern. We find some dependence on the used mass formula and on the distribution of the nuclear fragments after fissioning. Nevertheless, the second and third r-process peaks are robustly reproduced; without any ‘tuning’ of the nuclear physics input the overall agreement with the Solar system r-process pattern is good. r-process matter lighter than $\sim A \approx 120$, however, is substantially underproduced with respect to the Solar system pattern. The few high-$Y_e$ trajectories produce different abundance
underproduction of the lighter r-process elements, make compact robust abundance patterns for the heaviest r-process elements, but of a few, say, due to EOS or GR effects.

The typical ejecta amount per event should be decreased by a factor nicely brackets this value. This also implies that the relevance of the ejected masses are consistent with being a major contributor to the galactic r-process, even if the total amount per event should be decreased by a factor of a few (say, due to EOS or GR effects).

One may wonder whether captures of neutrinos from the central remnant, which are not included in our simulations, could change the $Y_e$ of the ejecta notably. For two reasons, we do not think this is the case. First, the ejecta are launched practically at first contact, see Fig. 2, with $(v) > 0.1c$ (Piran et al. 2012). The neutrino emission, however, only becomes relevant after a disc has formed which takes $\sim 7$ more milliseconds (e.g. fig. 7 in Rosswog et al. 2012) at which time the solid angle at which the escaping matter is seen has decreased already by a large factor. Secondly, the ejecta are concentrated in the orbital plane while the neutrinos are emitted preferentially perpendicular to it, see fig. 12 in Rosswog & Liebendörfer (2003) and fig. 9 in Dessart et al. (2009). We therefore expect $Y_e$-changes due to neutrino captures to be negligible.4

The dynamic ejecta masses are large enough to make a substantial contribution to the chemical enrichment of the galaxy with r-process material. Adopting the observation-based double ns merger rates of Kalogera et al. (2004), $R_{\text{DNS}} = 8.3^{+20.1}_{-6.6} \text{ yr}^{-1}$, we find out by which average rate the galaxy would be r-process enriched from ns mergers. Although the overall ejecta masses vary by more than a factor of 5 between different cases, all of them are perfectly consistent with the estimated galactic enrichment rate (Qian 2000), see Fig. 9. The reference case of $2 \times 1.4 \, M_\odot$ neutron stars (shown in magenta) nicely brackets this value. This also implies that the relevance of compact binary mergers for the r-process remains unaffected even if the typical ejecta amount per event should be decreased by a factor of a few, say, due to EOS or GR effects.

Our findings of ejecta masses in the required range, extremely robust abundance patterns for the heaviest r-process elements, but underproduction of the lighter r-process elements, make compact r-process sites must contribute to the cosmic chemical evolution of lighter r-process elements.

An important question is how early the r-process yields from compact binary mergers would be available for cosmic enrichment. The inspiral time from the birth of a compact binary system to a coalescence is very sensitive to both the initial separation and the orbital eccentricity (Peters & Mathews 1963, 1964). The initial separation is set by binary evolution processes which are beyond the scope of this paper, but it is worth re-iterating that a binary system that survives the two SN explosions that are required to form the system in the first place will in most cases possess a large orbital eccentricity. As illustrated in Fig. 10 short lifetimes can be achieved with plausible orbital parameters. For example, an initial semimajor axis of $1 \, R_\odot$ and an eccentricity of 0.9 lead to a merger only 1 Myr after binary formation [for comparison we note that the projected semimajor axis of the double pulsar PSR J0737$-$3039 A+B (Burgay et al. 2003; Kramer & Stairs 2008) is $0.6 \, R_\odot$]. Some binary population synthesis studies (Dewi & Pols 2003; Belczynski et al. 2006) find indeed short-lived channels that coalesce essentially at birth, so an early enrichment from compact mergers is at least plausible. Earlier studies (Argast et al. 2004) disfavoured ns mergers as a dominant galactic r-process source. Based on our hydrodynamic plus nucleosynthesis studies, however, we have shown that compact binary mergers are excellent production sites for the unique, heavy r-process component. We leave the question of whether or to which extent they are consistent with cosmic chemical evolution for future studies.

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