The multi-messenger picture of compact binary mergers

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In the last decade, enormous progress has been achieved in the understanding of the various facets of coalescing double neutron star and neutron black hole binary systems. One hopes that the mergers of such compact binaries can be routinely detected with the advanced versions of the ground-based gravitational wave detector facilities, maybe as early as in 2016. From the theoretical side, there has also been mounting evidence that compact binary mergers could be major sources of heavy elements and these ideas have gained recent observational support from the detection of an event that has been interpreted as a “macronova”, an electromagnetic transient powered by freshly produced, radioactively decaying heavy elements. In addition, compact binaries are the most plausible triggers of short gamma-ray bursts (sGRBs) and the last decade has witnessed the first detection of a sGRB afterglow and subsequent observations have delivered a wealth of information on the environments in which such bursts occur. To date, compact binary mergers can naturally explain most—though not all—of the observed sGRB properties. This article reviews major recent developments in various areas related to compact binary mergers.

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1. Introduction

In July 1974 Russel Hulse and Joseph Taylor discovered the pulsar PSR 1913+16 during a systematic search for new pulsars at the Arecibo Observatory in Puerto Rico. The pulsar was detected with a period of 59 ms, but with apparent changes of 80 µs from day to day. Soon it became clear that these changes are due to the Doppler shift caused by the orbital motion around an unseen companion star. The orbital period is only 7.75 hours and with a semi-major axes of about 3 $R_\odot$, but no observed eclipses, the companion star must also be a compact object, either a neutron star (ns) or a black hole (bh). The binary system is so tight that the orbital velocity is of order $10^{-3}c$ and general-relativistic effects build up quickly. To quantify the deviation of the dynamics from a purely Newtonian binary, one introduces so-
called “Post-Keplerian” (PK) parameters. Since these PK parameters each depend on both component masses, one can determine individual stellar masses through the measurements of at least two PK parameters and not just—as in Newtonian theory—the total binary mass via of Kepler’s third law. PSR 1913+16 has by now been observed for more than four decades and its parameters have been determined to an astonishing precision. For example, the current values\(^2\) for the neutron star masses are \(m_p = 1.4398 \pm 0.0002\, M_\odot\) for the pulsar and \(m_c = 1.3886 \pm 0.0002\, M_\odot\) for the companion star. The binary is a “clean” system in the sense that the components can be very accurately described as point masses and in particular tidal effects can be safely neglected. The periastron advances at a rate\(^2\) of \(\dot{\omega} = 4.226598 \pm 0.000005\) deg/yr (as compared to 43 arcsec per century for Mercury) and the orbit decays at a rate that agrees with General Relativity’s prediction for the emission of gravitational waves\(^3\)

\[
\dot{P}_b = -\frac{192\pi G^{5/3}}{5c^5} \frac{m_pm_c}{(m_p + m_c)^{1/3}} \left(\frac{2\pi}{P_b}\right)^{5/3} \times \frac{1 + (73/24)e^2 + (37/96)e^4}{(1 - e^2)^{7/2}}
\]

(1)

to about 0.3%.\(^2\) The latter, however, is only true if the relative Galactic acceleration difference between the pulsar and the solar system— they have a separation\(^2\) of about \((9.9 \pm 3.3)\) kpc—is taken into account. General Relativity had been probed in a number of solar system tests such as the perihelion shift of Mercury,\(^4\) gravitational light deflection\(^5\) or the Shapiro delay\(^7\) or, more recently the Lense-Thirring effect in satellite orbits\(^8\) but neither had the predicted radiative properties of gravity been probed nor regimes where the deviation from a flat space is substantially larger than in the solar system (gravitational potential \(\Phi/c^2 \sim 10^{-6}\)). The orbital decay of PSR 1913+16 was the first—though still indirect—confirmation of the existence of gravitational waves and it laid to rest a longstanding debate\(^10\) about the existence of gravitational waves and its quadrupolar nature. In 1993 Hulse and Taylor were awarded the Nobel Prize in Physics for their discovery of PSR 1913+16.

In 2003 an even better suited laboratory for strong field gravity\(^12\) was discovered\(^14\) the “double pulsar” PSR J0737-3039 where both components are active pulsars, usually referred to as pulsar A and B. With an orbital period of only 147 minutes, relativistic effects are even larger and with a distance of only about 1 kpc the systematic errors due to the relative Galactic acceleration are much smaller. The double pulsar has an eccentricity of \(e = 0.088\) and shows with \(\dot{\omega} = 16.8991 \pm 0.0001\) deg/yr an even larger periastron advance than PSR 1913+16. Its masses have again been very accurately measured as \(m_A = (1.3381 \pm 0.0007\) M_\odot\) and \(m_B = (1.2489 \pm 0.0007)\) M_\odot\) and within less than 10 years this binary has become the best testbed to probe the accuracy of gravitational quadrupole emission with a current accuracy well below 0.1%.\(^16\)

To date over 2000 pulsars are known, about 10% of them possess a binary companion\(^17\) and 10 systems consist of two neutron stars\(^18\) Such systems have turned out to be precious laboratories to probe gravitational theories\(^19\). Although binary evolution suggests that it would be natural to form also a reasonable fraction of
neutron star black hole binaries\cite{22,23} to date no such system has been found. They should, in a number of respects, have properties similar to double neutron star binaries and we will therefore follow the common practice to collectively refer to both types as “compact binary systems”.

Typically, one distinguishes gravity tests in different regimes: 1) quasi-stationary weak field, i.e. velocities $v \ll c$ and the spacetime is close to Minkowskian everywhere, 2) quasi-stationary strong-field, i.e. velocities $v \ll c$, but the spacetime can show substantial deviations from Minkowskian, 3) highly dynamical strong-field, i.e. $v$ becoming comparable to $c$ and involving a strongly curved spacetime and finally the 4) radiation regime where the radiative properties of gravity are tested. The first regime includes solar system tests, the second and fourth regime can be probed by well-separated neutron star binaries while regime 3) will be probed by the ground-based gravitational wave detectors such as LIGO, VIRGO and KAGRA\cite{29,30}.

They are expected to see the last inspiral stages ($\sim$ minutes) that are initially well described by post-Newtonian methods\cite{31,32} and the subsequent merger and ringdown phase where strong-field gravity, hydrodynamic and nuclear equation of state (EOS) effects from the neutron star matter determine the dynamics. For an overview over the numerical modeling of these phases we refer to recent reviews\cite{33,34}.

Compact binary mergers had also been suggested very early on as the production site for “rapid-neutron capture” or “r-process” elements,\cite{35,36} actually before the discovery of the Hulse-Taylor pulsar. Although it may seem a natural idea to form r-process elements in the decompression of the extremely neutron-rich neutron star material, this idea has for a long time been considered as somewhat exotic. This perception, however, has changed during the last one and a half decades, mainly because compact binary mergers (CBM) were found to robustly produce r-process elements without fine-tuning while core-collapse supernovae seem seriously challenged to providing the conditions for at least the heaviest r-process elements. This issue will be discussed in more detail in Sec. 4. The idea has been further boosted by the recent discovery of an nIR transient in the aftermath of the short GRB 130603B. Such transients from radioactive decays of freshly produced r-process elements had been a prediction of the compact binary merger model. In particular the characteristic time scale of $\sim$ one week and a peak in the nIR are consistent with extremely heavy elements with large opacities having been produced. This topic will be discussed in more detail in Sec. 4.4.

The merger of a compact binary system is also the most likely trigger of short gamma-ray burst (sGRBs). The idea that the sGRB phenomenon is caused by compact binary mergers has now stood the test of three decades of observation and it is considered the most plausible model for the “engine” behind sGRBs, although it is not completely free of tension with some observations. We will review in Sec. 5 the major arguments that link sGRBs with the idea of an compact binary merger origin, we will discuss how the predictions from the compact binary merger model compare to the observed GRB properties and where open questions remain.
In the following, various facets of compact binary mergers are discussed in more detail. We begin by discussing relevant time scales in Sec. 2, followed by the expected neutrino emission in Sec. 3, nucleosynthesis and closely related issues are treated in Sec. 4 and we discuss short GRBs in Sec. 5. We will conclude with a summary in Sec. 6.

2. Time scales

As a stellar binary system revolves around its centre of mass it emits energy and angular momentum via gravitational waves, which are both extracted from the orbital motion \(a\) with a power \(P \propto \omega_{\text{orb}}^{10/3}\). Since the total, kinetic plus potential, orbital energy is \(E_{\text{orb}} = -\frac{Gm_1m_2}{2a}\), a loss of energy means a reduction of the mutual separation, which, due to \(\omega_{\text{orb}} = \sqrt{GM/a^3}\), corresponds to an increase of the orbital frequency and therefore to a further enhanced gravitational wave emission. This runaway process finally leads to a merger of the two stars. For elliptical orbits, not only the semi-major axis \(a\), but also the eccentricity \(e\) evolves in time. For a non-zero eccentricity, its temporal change is negative, \(de/dt < 0\), and the eccentricity is “radiated away” very efficiently. The Hulse-Taylor pulsar, for example, has currently an eccentricity of \(e = 0.617\) and a semi-major axis \(a = 2.2 \times 10^9\) cm. By the time the orbit has shrunk to \(10\ R_{\text{ns}}\), however, it will have an eccentricity of \(e < 10^{-5}\), i.e. the orbit will be very close circular.

A non-zero orbital eccentricity can substantially shorten the inspiral time until coalescence. In general, the evolution equations for \(a\) and \(e\) need to be integrated numerically, but a good approximation for not too large eccentricities \([3]\) is given by

\[
\tau_{\text{GW}} \simeq 9.83 \times 10^6\text{years} \left(\frac{P}{\text{hr}}\right)^{8/3} \left(\frac{M}{M_\odot}\right)^{-2/3} \left(\frac{\mu}{M_\odot}\right)^{-1} (1 - e^2)^{7/2}. \tag{2}
\]

Contours of the inspiral time \(\log(\tau_{\text{GW}})\) for an equal mass binary system with \(1.4\ M_\odot\) per star are shown in Fig. 1. The inset shows the inspiral time (in seconds) for the last stages of a circular binary system as a function of the separation \(a_0\). Also marked are the locations of PSR 1913+16 and of the double pulsar PSR J0737-3039 although their masses are slightly different from \(1.4\ M_\odot\). Circular binaries can only merge within the lifetime of galaxy (\(10^{10}\) yrs) if they possess an initial orbital period below 16 hrs. The Hulse-Taylor pulsar has a life time of \(3 \times 10^8\) years until coalescence. A binary system with the same properties, but an initial eccentricity of \(e = 0.99\), in contrast, would merge within only \(10^4\) years. This dependence of the inspiral duration on eccentricity also has important implications for compact binary mergers as possible sources of r-process elements, see Sec. 4.3, and for their

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[a] Here, we will assume that we are dealing with a binary system of two non-spinning point masses in quadrupole approximation as first worked out by Peters and Mathews in 1963.[21]

[b] This estimate ignores an additional eccentricity-dependent factor which is to within a few percent equal to unity for \(e < 0.6\), and approaches monotonically \(\approx 1.8\) in the limit of \(e \to 1\), so that the estimates are always substantially better than a factor of 2.
Fig. 1. Inspiral time for a binary with $2 \times 1.4 \, M_\odot$. Shown are contours of $\log(\tau_{GW})$ (in years) as function of the current orbital period (in hours) and eccentricity. The filled circles indicate period and eccentricity of the Hulse-Taylor Pulsar PSR 1913+16 and the Double Pulsar J0737-039 (although their masses are not exactly 1.4 $M_\odot$). The inset focuses on the last inspiral stages where the orbits are to very high accuracy circular (note that the axes are different from the main plot). The last stages of the inspiral ($a_0 < 100$ km) only take fractions of a second.

The frequency of the emitted gravitational waves $f_{GW}$ as a function of the time until coalescence $\tau_{GW}$ is

$$f_{GW} = 134\text{Hz} \left( \frac{1.21M_\odot}{M_{\text{chirp}}} \right)^{5/8} \left( \frac{1\text{s}}{\tau_{GW}} \right)^{3/8},$$

(3)
where the “chirp mass” is $M_{\text{chirp}} = \mu^{5/3} M^{2/5}$. Therefore, a $M_{\text{chirp}} = 1.21 M_\odot$ binary would emit gravitational waves of 10 Hz or more (roughly the lowest frequencies accessible to ground-based detectors) during the last 17 minutes and 100 Hz or more during the last two seconds of the inspiral. For a double neutron star system, one expects the final merger to occur around gravitational wave frequencies of 1 kHz. The inspiral can be accurately described by means of Post-Newtonian methods,\textsuperscript{33,34} but the final merger and “ringdown” require numerical simulations.\textsuperscript{35–37} Simulations indicate that the central object that forms in a merger does in many cases not collapse directly to a black hole, but instead survives as a “hypermassive neutron star” (HMNS),\textsuperscript{44–47} despite having a mass in excess of the Tolman-Oppenheimer-Volkoff maximum mass. This is due to effects such as thermal pressure or –more importantly– differential rotation.\textsuperscript{48} A typical velocity field inside the merging binary system is shown in Fig. 2: at contact a shear interface forms for irrotational binaries

\textsuperscript{6} *Differential* rotation is much more efficient in stabilizing stars than uniform rotation since—in the center—the rotation can be very rapid without shedding mass at the surface.

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**Fig. 2.** Velocity fields during the merger of two initially non-spinning 1.4 M_\odot neutron stars.\textsuperscript{48} The shear interface, see panel one, becomes Kelvin-Helmholtz unstable and forms a string of vortex rolls. Subsequently the vortex rolls merge and the remnant finally forms a radially differentially rotating, so-called “hypermassive neutron star” that is temporarily stabilized against the gravitational collapse to a black hole.
(left panel) that becomes Kelvin-Helmholtz unstable (right panel) and evolves into a differentially rotating remnant. Differential rotation is very efficient in stabilizing stars\cite{44,49} and the time scale to collapse will be set by angular momentum transport processes that occur on a time scale that is substantially longer than the dynamical time scale

$$\tau_{\text{rot}} = 0.2 \text{ ms} \left(\frac{\rho}{\rho_{\text{nuc}}}\right)^{-1/2},$$

(4)

where $\rho_{\text{nuc}} = 2.65 \times 10^{14} \text{ g cm}^{-3}$ is the nuclear saturation density. Relativistic simulations of double neutron star mergers\cite{50,52} indicate the central remnant is a HMNS\cite{44} unless the initial ADM mass of the binary exceeds a threshold mass of 1.35 times the maximum mass of a cold, non-rotating neutron star. The discovery of two massive neutron stars, J1614-2230 with $M_{\text{ns}} = 1.97 \pm 0.04 M_\odot$\cite{53} and PSR J0348+0342 with $M_{\text{ns}} = 2.01 \pm 0.04 M_\odot$\cite{54} now place this limit to $M_{\text{thresh}} > 2.7 M_\odot$ so that —depending on the distribution of neutron star masses that is realized in binary systems— a large fraction of the mergers may go through such a metastable phase.

Taking the current mass estimates for binary neutron stars\cite{55} at face value (but keep in mind that some estimates have substantial errors) only three out of the 10 DNS systems could avoid such a metastable phase and would instead directly collapse to a black hole. The delay time between merger and bh formation also has a crucial impact on both the triggering of short GRBs and —since it substantially impacts on neutrino-driven winds— on the merger nucleosynthesis. These issues are further discussed in Sec. 4 and 5. A long-lived HMNS could also produce additional radiative signatures and contributions to high energy cosmic rays\cite{56,58}.

In most cases a CBM leads to the formation of a “thick” accretion disk where the scale height $H$ is comparable to the radius $R_{\text{disk}}$. The dynamical time scale of such a disk is roughly

$$\tau_{\text{dyn, disk}} \sim \frac{2\pi}{\omega_K} \approx 0.01 \text{ s} \left(\frac{M}{2.5 M_\odot}\right)^{-1/2} \left(\frac{R_{\text{disk}}}{100 \text{ km}}\right)^{3/2},$$

(5)

while its viscous time scale is

$$\tau_{\text{visc}} \sim 0.3 \text{ s} \left(\frac{0.05}{\alpha}\right) \left(\frac{R_{\text{disk}}}{100 \text{ km}}\right)^{3/2} \left(\frac{2.5 M_\odot}{M}\right)^{1/2} \left(\frac{R_{\text{disk}}/H}{3}\right)^2,$$

(6)

where we have assumed that the viscosity can be parametrized as a Shakura-Sunyaev-type dissipation\cite{59}.

\footnote{In the following, we will use the abbreviation HMNS collectively for both “hypermassive neutron stars” that will finally collapse into a black hole and also for very massive, but stable, neutron stars.}

\footnote{We do not distinguish here between “disk” and “torus” and use both words synonymously.}
3. Neutrino emission

Like supernovae, compact binary mergers emit a large fraction of the released gravitational binding energy ($\sim 10^{53}$ erg) in the form of neutrinos. This results predominantly in neutrinos in the energy range of $\sim 20$ MeV. Their moderate energies together with the steep energy dependence of the interaction cross sections make them hard to detect\textsuperscript{4}. But if the merger also accelerates material to relativistic speeds, as expected if they are the sources of short GRBs (Sec. 5), one may also expect neutrinos of substantially larger energy. The same shocks that are thought to accelerate the electrons responsible for the prompt gamma-ray emission should also produce relativistic protons, which can produce high-energy neutrinos\textsuperscript{5,6,9,71}

Internal shocks that can produce the observed $\sim$ MeV photons, should also be able to produce so-called prompt neutrino emission with energies in excess of $10^{14}$ eV\textsuperscript{65,68}. High-energy neutrinos may also be produced via neutron-rich outflow\textsuperscript{69,70} and reverse shocks\textsuperscript{71}. Such neutrino production mechanisms, however, rely on GRB outflows having a sizeable baryonic component. If instead the outflow should be essentially Poynting flux, large neutrino fluxes would be unexpected, unless they are produced efficiently in the forward shock\textsuperscript{72,73}.

Our main focus here, however, is the lower neutrino energy channel coming directly from the merger remnant. Until merger, the neutron stars are still in cold $\beta$-equilibrium, since tidal dissipation causes only a raise to moderate temperatures\textsuperscript{74} ($T \sim 10^8$ K) and the tidal interaction has practically no time to change the neutron-to-proton ratio by weak interactions, see the inset in Fig. 1 for an illustration. The $\beta$-equilibrium condition on the chemical potentials is given as\textsuperscript{5} $\bar{\mu}_n = \bar{\mu}_p + \bar{\mu}_e + \bar{\mu}_{\bar{\nu}_e}$, where the bar indicates that the particle mass contributions are included. Since in old neutron stars no neutrinos are present (their diffusion time is of order seconds, see below) the neutrino chemical potential vanishes and the difference between the nucleon chemical potentials equals the electron chemical potential which, in turn, determines the electron fraction $Y_e$ which has typical values below 0.1. After the merger, we can expect to have a fraction of the virial temperature, $T_{\text{vir}} \sim 25$ MeV ($M/2.5$ $M_\odot$) (100 km/$R$), i.e. temperatures substantially beyond the electron-positron pair production threshold ($m_e c^2 = 0.511$ MeV). Therefore, positron captures $n + e^+ \rightarrow p + \bar{\nu}_e$ drive the electron fraction to higher values and yield the copious emission of electron-type anti-neutrinos which are the dominant neutrino species. This is different from the core-collapse SN case that yields neutrino luminosities in a similar regime, but dominated by the $\nu_e$ from the neutronization process $p + e \rightarrow n + \nu_e$. In the hot, high-density interior of a HMNS one also expects heavy lepton neutrinos to be produced\textsuperscript{76,79}. Since they can only

\textsuperscript{4}Keep in mind that so far neutrinos have been detected only once for a supernova\textsuperscript{60} SN1987A, and compact binary mergers occur $\sim 10^3$ times less frequently.

\textsuperscript{5}We discuss here only the simplest case where neutrons, protons and electrons are present in the neutron star. When the Fermi energies become large enough to allow for new particle species to appear, the condition has to be modified accordingly, see for example Glendenning\textsuperscript{75}.
interact via neutral current reactions, they escape easier from hotter and denser regions and therefore possess higher energies. Clearly, for the conditions prevailing in the remnant ($\rho > 10^{14} \text{ gcm}^{-3}$ in the HMNS and \sim 10^{10} \text{ gcm}^{-3}$ in the disk) photons are completely trapped on the relevant time scales and neutrinos are the only cooling agents.

In the following, we will assume that a HMNS with temperature $T_{\text{HMNS}} = 20$ MeV, a mass $M = 2.5\, M_\odot$ and with a radius $R = 20$ km is present, so that its average density is $\bar{\rho} \approx 1.5 \times 10^{14} \text{ gcm}^{-3}(M/2.5\, M_\odot)(R/20\, \text{km})^{-3}$. We further assume for our estimates the presence of a surrounding accretion disk of $T_{\text{disk}} = 5$ MeV, $M_{\text{disk}} = 0.2\, M_\odot$, a typical radius $R_{\text{disk}} = 100$ km, a typical density $\rho_{\text{disk}} = 10^{11} \text{ gcm}^{-3}$ and an aspect ratio $H/R = 1/3$, where $H$ is the characteristic disk height.

To obtain an order of magnitude estimate for the neutrino properties, we assume for simplicity that the main source of opacity is provided by neutrinos scattering off nucleons $h$ with a cross-section given by $\sigma = (1/4)\sigma_0(E_\nu/m_e c^2)^2$, where the reference cross-section is $\sigma_0 = 1.76 \times 10^{-44} \text{ cm}^2$. Therefore, the mean free path in the HMNS is

$$\lambda_{\nu}^{\text{HMNS}} = \frac{1}{n\sigma} \approx 1.8 \text{ m} \left(\frac{R}{20 \text{ km}}\right)^3 \left(\frac{M}{2.5\, M_\odot}\right)^{-1} \left(\frac{E_\nu}{60 \text{ MeV}}\right)^{-2}$$

and the corresponding neutrino optical depth is

$$\tau_{\nu}^{\text{HMNS}} \sim \frac{R}{\lambda_{\nu}^{\text{HMNS}}} \sim 10^4 \left(\frac{R}{20 \text{ km}}\right)^{-2} \left(\frac{M}{2.5\, M_\odot}\right) \left(\frac{E_\nu}{60 \text{ MeV}}\right)^2$$

i.e. the HMNS is very opaque to its own neutrinos. The characteristic neutrino diffusion time scale is then

$$\tau_{\nu}^{\text{diff}} \sim 3\tau_{\nu}^{\text{HMNS}} \frac{R}{c} \sim 2.4 \text{ s} \left(\frac{R}{20 \text{ km}}\right)^{-1} \left(\frac{M}{2.5\, M_\odot}\right) \left(\frac{E_\nu}{60 \text{ MeV}}\right)^2,$$

i.e. on the HMNS dynamical time scale, see Eq. (4), the neutrinos are efficiently “trapped”. If we assume that a thermal energy of $\Delta E_{\text{th}} \sim 0.1 GM^2/R$ can in principle be emitted, one expects a HMNS luminosity of

$$L_{\nu}^{\text{HMNS}} \sim \frac{\Delta E_{\text{th}}}{\tau_{\nu}^{\text{diff}}} \sim 3 \times 10^{52} \text{ erg/s} \left(\frac{M}{2.5\, M_\odot}\right) \left(\frac{E_\nu}{60 \text{ MeV}}\right)^2,$$

and, applying similar estimates as before, one finds the mean free path in the disk

$$\lambda_{\nu}^{\text{disk}} \approx 42 \text{ km} \left(\frac{\rho}{10^{11} \text{ gcm}^{-3}}\right)^{-1} \left(\frac{E_\nu}{15 \text{ MeV}}\right)^{-2}$$

\footnotetext[1]{For $\nu_e$ the opacity related to the absorption by neutrons is even larger, but still of the same order.}
\footnotetext[2]{The expected neutrino energies are moderate multiples of the matter temperature, usually given by the ratio of two Fermi-integrals at the local chemical potential. For our scalings we use $E_\nu \approx 3 k_B T$, i.e. 60 MeV for the HMNS and 15 MeV for the disk.}
\footnotetext[3]{We abbreviate both the optical depth and typical time scales with the symbol $\tau$. Since each time it has unique subscript this should no lead to confusion.}
and therefore the optical depth is $\tau \sim H/\lambda_{\nu}^{\text{disk}} \sim 1$. Thus, the typical escape time for a neutrino is substantially shorter than the dynamical and viscous time scales of the disk, see Eqs. (5) and (6). This short cooling time scale implies that neutrino emission can only be maintained through the constant supply of thermal energy via accretion. The gravitational energy gained by the accretion of the disk is then $\Delta E \sim GM_{\text{disk}}/R$ and the expected neutrino luminosity becomes

$$L_{\nu}^{\text{disk}} \sim \frac{1}{2} \frac{\Delta E}{\tau_{\text{visc}}} \sim 10^{53} \text{erg} \left( \frac{M}{2.5M_{\odot}} \right)^{\frac{3}{2}} \left( \frac{M_{\text{disk}}}{0.2M_{\odot}} \right) \left( \frac{\alpha}{0.05} \right) \left( \frac{R_{\text{disk}}}{3H} \right)^{2} \left( \frac{100 \text{ km}}{R_{\text{disk}}} \right)^{\frac{3}{2}} \left( \frac{20 \text{ km}}{R} \right)^{2}. $$

Clearly, these are only order of magnitude estimates, for more precise results numerical simulations are required. To date only a few implementations of neutrino physics in merger simulation codes exist.$^{76,77,81-83}$ It is encouraging that despite the steep temperature dependence of the neutrino emission rates (the electron capture energy emission rates, for example,$^{84}$ are $\propto T^{6}$), the predicted neutrino luminosities$^{77,83,85-87}$ are rather robust, for a $2 \times 10^{53}$ erg/s with average energies around 10, 15 and 20 MeV for $\nu_e$, $\bar{\nu}_e$ and $\nu_x$, respectively, and, as expected from the simple considerations, with the luminosity being dominated by $\bar{\nu}_e$. It seems in particular that implementation details play less of a role than the physical ingredients such as the treatment of gravity or the equation of state.

While the quoted results are based on relatively simple leakage schemes, more sophisticated neutrino treatments$^{78,79}$ yield actually similar results. It had however, been realized by Dessart et al.$^{78}$ that nucleon-nucleon bremsstrahlung, which had been ignored in early implementations, is also an important source of heavy lepton neutrinos.

More recently, neutrino emission has also been included in nsbh-merger simulations.$^{82,88}$ Here the parameter range, especially the expected mass ratio range, is in principle even larger (and to date not yet well explored), but the first studies$^{88}$ find, for $M_{\text{bh}} = 7 M_{\odot}$ and large bh spin ($a_{\text{bh}} > 0.7$), also luminosities around $\sim 2 \times 10^{53}$ erg/s. Lower mass black holes are expected to yield hotter disks and therefore larger neutrino luminosities. Indeed, a simulation for a 5.6 $M_{\odot}$ bh with $a_{\text{bh}} = 0.9$ yield peak luminosities $> 10^{54}$ erg/s. The bh mass distribution in nsbh binaries, however, is not well known and a recent study$^{89}$ based on X-ray binaries seems to suggest a peak around $\sim 8 M_{\odot}$. These results are consistent with studies$^{90}$ that find, from a binary evolution point of view, bh masses near $\sim 10 M_{\odot}$ most likely.

4. Compact binary mergers as cosmic factories of the heaviest elements

The heaviest nuclei in the Universe form via neutron captures. In this process there is a competition between captures of new neutrons and $\beta-$decays and the ratio of the two time scales can be used to distinguish between a “slow” (s-process) and
a “rapid neutron capture process” (r-process). Each process is responsible for about half of the elements heavier than iron. Although the physical mechanisms have been understood decades ago, there is still no consensus about the astrophysical production site of the r-process. For many years supernovae have been considered as the most promising site, in particular the neutrino-driven wind that emerges after the formation of a hot, still deleptonizing proto-neutron star. Over the last decade, however, via increasingly sophisticated supernova simulation models, the consensus has grown that the neutrino-driven winds do not provide the right physical conditions for the production of the third r-process peak \((A \sim 195)\), but contributions to the less heavy region up to the second peak near \(A \sim 130\) may be plausibly expected, see Arcones and Thielemann for an excellent review. A possible alternative supernovae site could be magnetohydrodynamically launched jets, but these require a rather extreme combination of spin and magnetic field as initial conditions and at present it is not clear whether or at which rate such conditions are realized in nature.

Neutron-rich matter ejected in a compact binary merger has been considered as an alternative formation channel of rapid neutron capture (“r-process”) elements since the 70ies. Despite its physical plausibility, this channel was considered exotic for a long time and only the second-best model after core-collapse supernovae. This perception has only changed in the last one and a half decades through a number of studies that found r-process, up to the third r-process peak \((A \sim 195)\), naturally occurring in compact binary merger ejecta. The ease with which r-process is produced by compact binary mergers together with the serious difficulties to produce heavy r-process in core-collapse supernovae has lead to a shift in the general opinion towards compact binary mergers as the most likely production site of at least the heaviest, but possibly of even all of the r-process elements. An important issue that needs to be better understood, however, is whether/how r-process from compact binary mergers is consistent with galactochemical evolution.

4.1. Mass loss channels

During the merger of a compact binary system mass is ejected into space via a number of channels. A fraction is released dynamically, i.e. via hydrodynamic interaction and gravitational torques, and will subsequently be referred to as “dynamic ejecta”. Most CBMs also result in an accretion disk, and as such a disk evolves on a viscous time scale, a substantial fraction of its initial mass becomes unbound as a result of nuclear and viscous action. Another channel that has received a fair amount of attention recently are neutrino-driven winds. Highly magnetized neutron star-like merger remnants can also magnetically drive winds. Also high-velocity ejecta coming from the interaction region between two neutron stars have been postulated. All these channels differ in the amount of ejected mass, the distribution of electron fractions and entropies. There-
fore, they produce likely different element distributions and –if all produce electromagnetic transients– these may be different for each channel.

What is an ejecta amount that is interesting from a galactic r-process enrichment perspective? If $R_i$ denotes the rate of event type $i$ (averaged over the Galactic age) and $\bar{m}_{i,c}$ is the average ejected mass per event in a particular channel $c$ of this event type, then the average Galactic r-process enrichment rate is

$$\dot{M}_{r,\text{gal}} = \frac{M_{r,\text{gal}}}{\tau_{\text{gal}}} = \sum_i R_i \sum_c \bar{m}_{i,c} = \sum_i R_i \bar{m}_i,$$

where $\dot{M}_{r,\text{gal}}$ is the r-process enrichment rate, averaged over the age of the Galaxy $\tau_{\text{gal}}$. In the following estimates we will conservatively restrict our considerations to nucleon numbers $A > 140$, since a large number of studies has shown that dynamic ejecta robustly produce such heavy r-process elements. It is, however, also a possibility that CBM produce all r-process material and it is also possible that event types different from CBMs contribute to the r-process production. This latter possibility will be ignored in our below estimates, therefore the quoted numbers should be interpreted as upper limits.

If we assume that the solar abundance pattern is representative for the Milky Way, we can multiply the mass fraction of nuclei with $A > 140$, $X_r^{A>140} = 2.6 \times 10^{-8}$, with the baryonic mass of the Milky Way $6 \times 10^{10} M_\odot$, to find a Galactic r-process mass ($A > 140$) $M_{r,\text{gal}}^{A>140} \approx 1560 M_\odot$ and therefore the average r-process enrichment rate ($A > 140$) is $\dot{M}_{r,\text{gal}}^{A>140} \approx 1.1 \times 10^{-7} M_\odot \text{ year}^{-1}$. Since

Fig. 3. A schematic illustration of the merger remnant. The numerical values are from recent numerical studies.

\[Y_e \approx 0.31 - 0.35\]
\[s \approx 15 - 20\text{ km/bar.}\]
\[v \approx 0.09 \text{ c}\]

\[Y_e \approx 0.23 - 0.31\]
\[s \approx 15 \text{ km/bar.}\]
\[v \approx 0.06 \text{ c}\]

\[Y_e \approx 0.04\]
\[v \approx 0.10 - 0.18 \text{ c}\]
nsbh mergers are estimated to be at least an order of magnitude rarer than nsns mergers\footnote{Using simulation results on ejected mass\textsuperscript{115} the nsbh rate has recently been constrained to $R_{\text{nsbh}} < 60 \text{MWEG}^{-1}\text{Myr}^{-1}$.}.\footnote{For an example see Fig. 2 in Korobkin et al. 2012\textsuperscript{115}} We only consider the latter for the following order of magnitude estimate\footnote{Using "plausible pessimistic" and "plausible optimistic" rates\textsuperscript{136} as brackets, $R_{\text{nsns}}^{\text{low}} \sim 1 \text{MWEG}^{-1}\text{Myr}^{-1}$ and $R_{\text{nsns}}^{\text{high}} \sim 1000 \text{MWEG}^{-1}\text{Myr}^{-1}$, where MWEG refers to "Milky Way-equivalent Galaxy", we see that \begin{equation}
\bar{m}_{\text{nsns}} \sim \dot{\mathcal{M}}_{r,\text{gal}}/R_{\text{nsns}} \sim 10^{-4} ... 0.1 \ M_{\odot}
\end{equation}}

\begin{equation}
\bar{m}_{\text{nsns}} \sim \dot{\mathcal{M}}_{r,\text{gal}}/R_{\text{nsns}} \sim 10^{-4} ... 0.1 \ M_{\odot}
\end{equation}

is an interesting ejecta amount from a chemical evolution perspective. If instead of only nucleon numbers $A > 140$ all r-process elements should be produced in CBM, the ejecta masses would need to be larger by a factor of $X_r/X_r^{140} \approx 13$, if there are substantial contributions from other events the nsns yields need to be correspondingly lower. These estimates are what hydrodynamic simulation results need to be compared with. It further needs to be understood whether such masses/rates are consistent with the observed chemogalactic evolution.

4.1.1. Dynamic ejecta

In fact, essentially all recent binary merger calculations find dynamic ejecta masses in a range that is consistent with being a major source of cosmic r-process, see Table \ref{tab:dynamic_ejecta}.

\begin{table}[h]
\centering
\begin{tabular}{lcccc}
\hline
authors & reference & dyn. ejecta [0.01 $M_{\odot}$] & binary & comment \\
\hline
Bauswein et al. 2013 & \textsuperscript{135} & $1.7 \times 10^{-2}$ ... 1.8 & nsns & CF-approximation \\
Bauswein et al. 2014 & \textsuperscript{137} & $< 2 \times 10^{-4}$ ... 9.6 & nsbh & CF-approximation \\
Deaton et al. 2013 & \textsuperscript{79} & 8 & nsbh & $M_{bh} = 5.6 \ M_{\odot}$, $a_{bh} = 0.9$ \\
Foucart et al. 2013 & \textsuperscript{139} & 1 ... 5 & nsbh & $M_{bh} = 10 \ M_{\odot}$, $a_{bh} = 0.9$ \\
Foucart et al. 2014 & \textsuperscript{86} & 5 ... 20 & nsbh & $M_{bh} = 7...10 \ M_{\odot}$, $a_{bh} = 0.7...0.9$ \\
Hotokezaka et al. 2013 & \textsuperscript{140} & 0.01 ... 1.4 & nsns & full GR \\
Kyutoku et al. 2013 & \textsuperscript{114} & 1 ... 7 & nsbh & full GR \\
Oechslin et al. 2007 & \textsuperscript{112} & 0.1 ... 4.5 & nsns & CF-approximation \\
Rosswog et al. 2013 & \textsuperscript{119} & 0.76 ... 3.9 & nsns & Newt. \\
Rosswog 2005 & \textsuperscript{113} & 1 ... 20 & nsbh & $M_{bh} \geq 14 \ M_{\odot}$, PW-potential \\
\hline
\end{tabular}
\caption{Comparison of the masses for the dynamic ejecta found in different numerical studies. While the numbers give a good impression about the expected mass range, it is worth keeping in mind that they are of limited comparability since the studies vary different parameters. CF: conformal flatness approximation, GR: General Relativity, PW: Paczynski-Wiita potential.}
\end{table}

The matter from the interaction region is\footnote{For an example see Fig. 2 in Korobkin et al. 2012\textsuperscript{115}} that emerges from the interface between the two neutron stars and a "tidal component"\footnote{For an example see Fig. 2 in Korobkin et al. 2012\textsuperscript{115}} that is launched by gravitational torques without experiencing noticeable shocks or shear flows.\textsuperscript{112,114}
considerably hotter than the tidal component. Therefore, it may, via positron captures, change its electron fraction $Y_e$, while the tidal component is ejected at its low, initial $\beta$-equilibrium value. The ratio of the mass in these components also depends on the dynamic evolution at merger and therefore on the treatment of gravity and the used equation of state. In Newtonian calculations\cite{87,115,118} with the stiff Shen et al. EOS\cite{143,144} the tidal component dominates while simulations using the Conformal Flatness approximation to GR\cite{138} and simulations using dynamical, general-relativistic space-times\cite{140} are dominated by the interaction component. Due to their possibly different electron fractions this may have implications for nucleosynthesis, see below.

4.1.2. Disk dissolution

When an accretion disk forms in the aftermath of a compact binary merger, it is characterized by $\rho \approx 10^{10}$ g cm$^{-3}$, $T \approx 10^{10}$ K and initial electron fractions close to the $\beta$-equilibrium value, $Y_e < 0.1$, of the disrupted neutron star, but increasing as the disk evolves viscously. Under such conditions the bulk of the disk consists of free neutrons and protons and $e^{+}e^{-}$-pairs. As the disk evolves viscously on a secular time scale, see Eq. (6), it expands and cools and at some point the free nucleons combine into light elements, predominantly alpha particles\cite{117} which releases $\approx 7$ MeV/nucleon. Together with viscous dissipation, turbulent energy transport and energy deposition by neutrinos this can unbind a substantial fraction (up to 25\%) of the original disk mass\cite{117,119,122,145}. How much matter becomes unbound by neutrino absorption is crucially dependent on the the presence of a central HMNS vs a black hole, in the latter case substantially less material is blown off from the disk\cite{78,79,117,122,123}.

4.1.3. Neutrino-driven winds

When two neutron stars merge, they release a substantial fraction of the gravitational binding energy in the form of neutrinos. As shown in Sec. 3 the remnant neutrino luminosities are in the range of a few times $10^{53}$ erg/s with typical neutrino energies of $\sim 15$ MeV\cite{76,77,83,86,146}. The gravitational binding energies per nucleon in the merged remnant, in comparison, are of order $E_{grav} \approx -35$ MeV ($M_{co}/2.5 M_\odot$) (100 km/$r$), where $M_{co}$ is the mass of the central object. That means capturing a few neutrinos provides enough energy to potentially escape the gravitational attraction of the remnant. It had been realised early\cite{16,124,125} that the absorption of neutrino energy could drive strong baryonic winds, similar to the case of a new-born proto-neutron star\cite{96,97} see Fig. 3 for a simple sketch, but the wind properties (mass, electron fraction, entropy or geometry) are difficult to estimate without numerical simulations.

\footnote{Keep in mind that on the viscous time scale weak interactions may have substantially shifted the electron fraction $Y_e$ to larger values.}
Early investigations used either order-of-magnitude approaches or parametrized models to explore the neutrino-winds from merger remnants\cite{111,119,123,126,147,151}. Due to the involved technical challenges, full-fledged numerical simulations of neutrino-driven winds have only been performed recently. The first neutrino-hydrodynamic study was performed in 2D by Dessart et al.\cite{78} followed by recent studies by Fernandez and Metzger\cite{123,145} and by Just et al.\cite{117} Just et al. explored the different ejecta channels using viscous hydrodynamics with Newtonian and pseudo-Newtonian gravity and they applied an energy-dependent two-moment closure scheme for the transport of electron neutrinos and anti-neutrinos. For black hole disk systems, they found neutrino interactions to help unbind disk material ($\sim 1\%$ of the initial disk mass), however, to a much smaller extent than viscous action which is able to unbind up to $25\%$ of the initial disk mass. This result is consistent with earlier studies in 2D\cite{123} that also found that neutrino heating in bh-disk systems unbinds only a small fraction of the disk material.

To date, we are aware of only a single study in 3D, performed by Perego et al.\cite{152} They explored the case where a HMNS is present and used the end points of 3D SPH simulations\cite{48} as initial condition. The further evolution was followed by means of an Eulerian hydrodynamics code\cite{153} augmented by a detailed, spectral leakage scheme that has been gauged at Boltzmann transport calculations of supernovae. They found a strong baryonic wind being blown out along the original binary rotation axis within $\sim 100$ ms, in qualitative agreement with the earlier 2D work by Dessart et al.\cite{78} This wind unbinds at least $3.5 \times 10^{-3}$ M$_\odot$ and is therefore of relevance for the galactic chemical evolution, see Eq. (14). The weak interactions in the wind yield a broad distribution of electron fractions between $0.2$ and $0.4$, with tendentially more proton-rich outflow ($Y_e > 0.3$) along the rotation axis while the equatorial outflow remained more neutron-rich ($0.2 < Y_e < 0.3$).

\section*{4.2. Nucleosynthesis in different channels}

All of the discussed mass loss channels are possible r-process sources. Most advanced are to date the calculations for the dynamic ejecta, the other channels are still in early exploration stages. Already the early network calculations for the dynamic ejecta of compact binary mergers\cite{111} showed a very good agreement with the observed solar system abundance pattern, see Fig. 4, left panel. Since no weak interactions were included in these calculations, the electron fraction $Y_e$ was varied in a low range considered appropriate for neutron stars. These results showed a very good agreement with the solar system pattern beyond $A \approx 130$ and in particular produced the heaviest elements up to the platinum peak without any fine tuning of the parameters, but hardly produced yields below $A = 130$. Such calculations have been more and more refined over recent years\cite{113,115,138,154,155} and they all agree that the dynamic ejecta are a very promising site for (at least) the heaviest elements ($A > 130$). The abundance pattern has been shown\cite{113} to be virtually independent of the astrophysical parameters of the merging binary system (masses,
mass ratio, whether the binary system consists of two neutron stars or a neutron star and a black hole) so that every CBM produces the same, very robust abundance pattern. This property is, of course, very interesting from a galactochemical point of view, since such robust abundance patterns are actually observed for the heaviest elements ($Z \geq 56$) in metal-poor stars\textsuperscript{133}. Moreover, studies that vary the neutron star EOS\textsuperscript{138} come to the conclusion that the abundance is also rather stable against such variations. A recent study\textsuperscript{155} has varied nuclear mass models and investigated the reason for the robustness. The authors suggest that a requirement for the robustness is that at the moment of freeze-out a much larger mass is in the fissioning region (nucleon number $A > 250$) than in the second r-process peak and above ($120 < A < 180$).

While the results are –for fixed nuclear physics ingredients– very robust with respect to a change of the astrophysical parameters, they show a substantial sensitivity with respect to the variation of the nuclear physics input, they are particularly sensitive to the distribution of the fission products\textsuperscript{115,155,156} which have a particular influence on the abundance pattern after the second peak. Although the abundance pattern globally agrees well with the solar system abundances, a close inspection shows in a number of studies\textsuperscript{111,113,115,138,157,159} a slight, but systematic shift (with respect to the solar pattern) of the calculated results towards higher $A$. These are attributed to “late” neutrons, e.g. from $\beta$-delayed neutron emission, but the agreement can be improved\textsuperscript{156,160} if $\beta$-decay rates are applied that are faster than those derived from the frequently used Finite Range Droplet Model\textsuperscript{161}. Such increased rates are actually consistent with recent experimental data for nuclei close to the r-process path\textsuperscript{162} and recent shell model calculations\textsuperscript{163,164}.

The nucleosynthesis in the other channels has so far been much less explored. In

![Fig. 4. Left panel: r-process abundance pattern found the dynamic ejecta from neutron star mergers\textsuperscript{111} are shown as line, the solar system abundance pattern as crosses. Right panel: sketch of different event locations with respect to their host galaxy.](image)

both the neutrino-driven wind and the disk-dissolution channel matter stays much
longer, roughly a viscous time scale Eq. (6) at larger temperatures and in a neutrino background field. Therefore, the weak interactions have time to substantially increase the electron fractions of the finally unbound material. The first hydrodynamic studies of these channels\textsuperscript{79,117} find that they produce less heavy r-process elements and with “non-robust” abundance patterns that depend on the exact details on the merged system. Again, such a behaviour is consistent with the observation of metal-poor stars\textsuperscript{134} that show substantial deviations from the solar system pattern for $Z < 56$ on a star-by-star basis. The wind- and disk-yields complement the dynamic ejecta at nucleon numbers below $A \approx 130$ and show broad consistency with the solar abundance pattern.

Relativistic gravity also seems to have a substantial influence on the resulting nucleosynthesis. This is mainly because neutron stars are more compact and the dynamics close to merger is impacted by GR effects. Overall, one expects a more violent dynamics, larger temperatures in the merged remnant and therefore an enhanced likelihood to achieve a broader distribution of $Y_e$ values. A recent study\textsuperscript{116} including general relativity and a detailed neutrino treatment reports, for example, on a reproduction of the whole r-process range, however, without discussing in detail the channels and mechanisms. Closely related to the impact of GR is the equation of state and the GR-effects can be enhanced by particularly soft equations of state, while stiffer ones yield results that are closer to those employing Newtonian gravity.

4.3. \textit{Galactic chemical evolution}

As outlined in Sec. \textsuperscript{12} all recent studies based on hydrodynamic source plus nuclear network simulations come to the conclusion that compact binary mergers are excellent candidates for a major r-process source, at the very least for the heaviest nuclei ($A > 130$), but maybe even for all r-process nuclei. But there are, of course, also constraints from galactochemical evolution and the observed elemental scatter in metal-poor stars should be different depending on whether the majority of r-process comes from core-collapse supernovae or from compact binary mergers. For example, as discussed in Sec. \textsuperscript{4} the typical ejecta mass of CBM of $\sim 0.01$ M$_\odot$ would, together with typical rate estimates, be able to account for most/all r-process material. If core-collapse supernovae, in contrast, would be main source and each supernovae would contribute, the average ejecta mass per event would only be $\sim 10^{-5}$ M$_\odot$. Moreover, as will be discussed in the sGRB context, see Sec. \textsuperscript{5} the large sensitivity of the gravitational wave inspiral times to the initial binary separation and eccentricity, see Eq. (2) and Fig. 1, naturally leads to a large spread in delay times after star formation. If endowed with substantial kick velocities, one expects a large variety of merger sites: some would merge close to their birth places in the galactic plane while others would merge in the outskirts of galaxies, see the

\textsuperscript{a}In cases where a stable neutron star results the time scales could be as long as the neutrino cooling time scales of several seconds, see Eq. (9).
right panel of Fig. 4 for a sketch. Such a distribution is indeed observed for short
GRBs (also thought to result from CBM), they show typical projected offsets of
\( \sim 5 \) kpc, see Sec. 5.1. This also implies a large range of ambient matter densities
encountered for different merger sites and therefore a different “stopping length”
and volume with which the ejecta can mix. Thus, the galactochemical imprint of
SN and CBM is likely quite different.

Initial inhomogeneous chemical evolution studies\cite{165} ruled out neutron star mergers
as dominant source of r-process, mainly since –due to the low rates– the r-process
enrichment would not be consistent with the observations at low metallicity and
the scatter in [r-process/Fe] would be much too large. This is influenced on the one
hand by the delay time due to GW inspiral, see Eq. (2) and Fig. 1, and on the other
hand by the question whether the neutron star ejecta mix with the chemical prod-
ucts from the supernovae which formed the neutron stars in the first place. There
are, however, indications that at least some of those supernovae that form neutron
stars in surviving double neutron star systems are “non-standard”, they may form
low-mass, low-kick neutron stars and eject only very small amounts of heavy ele-
ments\cite{168,169}. More recent chemical evolution studies\cite{170} come to the conclusion that
neutron star mergers could well be a major contributor to the galactic r-process in-
ventory, although probably not the only one. Another chemical evolution study\cite{171}
concludes that, except for the earliest evolutionary phases, CBM could well be the
major production sites of r-process elements.

Most recently, hydrodynamic studies of galactic enrichments have become available.
In particular a high resolution cosmological hydrodynamics simulation called “Eris”
has been used that is thought to closely represent the evolution of the Milky Way\cite{172}
The results of this simulation have been used in a post-processing step to follow
the evolution of r-process enrichment by either type II supernovae or compact bi-
nary mergers. Unlike in previous studies, the authors find that the nucleosynthetic
products from compact binary mergers can be incorporated into stars of very low
metallicity and at early times, even with a minimum delay time of 100 Myr and
come to the conclusion that compact binary mergers could be the dominant source
of r-process in the Galaxy. These results are supported by an independent hydro-
dynamic study\cite{173} that uses cosmological zoom-in simulations of a Milky Way-mass
galaxy from the Feedback In Realistic Environments project\cite{175}. Also here, the
authors conclude that the results are consistent with ns mergers being the source of
most of the r-process nuclei in the Universe.

Such an interpretation is also consistent with a recent study\cite{174} of radioactive \(^{244}\)Pu
in Earth’s deep-sea reservoirs. The authors find that the abundances are two or-
ers of magnitude lower than what is expected from a continuous enrichment by
supernovae. This points to a very rare actinide production site, for example a small
subset of actinide-producing supernovae or a compact binary merger origin.
4.4. “Macronovae”: electromagnetic transients powered by radioactive decays from freshly synthesized heavy elements

As outlined above, compact binary mergers eject, via several channels, initially extremely neutron-rich matter. The first and best-studied channel are the dynamic ejecta. Li and Paczynski were the first to realize that if indeed rapid neutron capture occurs in the ejected material, then the radioactivity in the ejecta should cause an electromagnetic transient following a merger. Such transients are often referred to as “kilonovae” or “macronovae” But since only little mass is ejected (\( m_{\text{ej}} \sim 0.01 \, \text{M}_\odot \)) at large velocities (\( v_{\text{ej}} \sim 0.1 \, \text{c} \)), the expected light curve evolution is faster and the peak luminosity is lower than in the type Ia case. The bulk of the energy release due to r-process occurs on a time scale of ~ second when the matter density is still so large that radiation is completely trapped and the radioactive heating is used up for accelerating the adiabatic expansion of the ejecta. It is only when the expansion time scale becomes comparable to the radiative diffusion time scale that substantial electromagnetic emission occurs. The diffusion time scale and therefore the time to transparency depend on how opaque the material is with respect to its radiation content, but unfortunately the opacities are to date maybe the largest uncertainty in modeling macronovae.

For simplicity, the first study of Li and Paczynski (and all subsequent studies until recently) assumed for a lack of better knowledge that the relevant opacities would be of order \( \kappa \sim 0.1 \, \text{cm}^2/\text{g} \), similar to the line opacities of iron-group elements. The composition of the dynamic ejecta, however, is very different from any other known cosmic explosion and in particular very different from any type of supernova. While the latter produce elements up to the iron group near \( Z = 26 \), the dynamic ejecta of neutron star mergers consist, due to the extreme initial neutron-richness of the decompressed r-process material, almost entirely of r-process elements up to the third peak near \( Z \approx 90 \), e.g. Fig. 4. To date, not much is known about the relevant opacities of such material, but it was recently realized by Kasen and collaborators that the opacities have previously been underestimated by orders of magnitude. The literally millions of lines, Doppler-broadened by the differential velocities of the ejecta, contribute a pseudo-continuum of bound-bound opacity. Photons trying to escape the ejecta come into resonance with multiple transitions one by one, resulting in a very large opacity. The bulk of the opacity is thereby provided by ions with a particularly complex structure. In particular lanthanides (\( 58 \leq Z \leq 72 \)) have –due to the complicated structure of their valence f shells– been found to be major sources of opacity. While detailed knowledge of opacities is still lacking, the currently existing calculations indicate opacities that are \( \sim \) two orders of magnitude larger than what was previously assumed, \( \kappa \sim 10 \, \text{cm}^2/\text{g} \), even if only small amounts of lanthanides are present. As a result, radiation is much longer trapped and matter becomes transparent only much later, at lower temperatures.

\( ^{\text{a}} \text{For the nuclear heating history see, for example, Korobkin et al. their Fig.7.} \)
and luminosities and peaking in the near-IR band rather than, as initially found, in the optical. Using such increased opacities, recent studies \cite{182,183,184} found the resulting “macronovae” should peak after typically a week rather than the original estimate of half a day.

Simple order of magnitude estimates \cite{157,176,184} can be obtained in the following way. The diffusion time scale for the ejecta with radius $R$ is approximately

$$\tau_{\text{diff}} \sim \frac{m_{\text{ej}} \kappa}{4 \pi c R}.$$  \hspace{1cm} (15)

The peak emission is expected \cite{179} when this is equal to the expansion time scale $\tau_{\text{exp}} = \frac{R}{v_{\text{ej}}}$, so at a radius

$$R_{\text{peak}} \sim \sqrt{\frac{m_{\text{ej}} \kappa v_{\text{ej}}}{4 \pi c}} = 1.3 \times 10^{15} \text{ cm} \left( \frac{m_{\text{ej}}}{10^{-2} M_\odot} \right)^{1/2} \left( \frac{\kappa}{10 \text{ cm}^2/\text{g}} \right)^{1/2} \left( \frac{0.1 \text{c}}{v_{\text{ej}}} \right)^{1/2},$$  \hspace{1cm} (16)

which is reached after

$$\tau_{\text{peak}} \sim \sqrt{\frac{m_{\text{ej}} \kappa v_{\text{ej}}}{4 \pi c}} = 4.9 \text{ days} \left( \frac{m_{\text{ej}}}{10^{-2} M_\odot} \right)^{1/2} \left( \frac{\kappa}{10 \text{ cm}^2/\text{g}} \right)^{1/2} \left( \frac{0.1 \text{c}}{v_{\text{ej}}} \right)^{1/2}. \hspace{1cm} (17)$$

If we assume that the late-time radioactive energy release rate can be approximated as a power law \cite{157,159,176}, $\dot{\epsilon} = \dot{\epsilon}_0 (t/t_0)^{-\alpha}$ ($\alpha = 1.3$), we can estimate the peak bolometric luminosity as

$$L_{\text{peak}} \sim m_{\text{ej}} \dot{\epsilon}(t_{\text{peak}}) = \dot{\epsilon}_0 (4 \pi c)^{\frac{2}{3}} \left( \frac{v_{\text{ej}}}{\kappa} \right)^{\frac{2}{3}} \left( \frac{10 \text{ cm}^2/\text{g}}{\kappa} \right)^{\frac{2}{3}} \left( \frac{0.1 \text{c}}{v_{\text{ej}}} \right)^{\frac{2}{3}} \left( \frac{10^{-2} M_\odot}{m_{\text{ej}}} \right)^{1-\frac{2}{8}}, \hspace{1cm} (18)$$

where we have used the numerical values as determined in Korobkin et al. (2012). Using the Stefan-Boltzmann law, $L = 4 \pi R^2 \sigma_{SB} T^4$, one can obtain a rough estimate for the effective temperature at peak as

$$T_{\text{peak}}^{\text{eff}} = 2200K \left( \frac{10 \text{ cm}^2/\text{g}}{\kappa} \right)^{\frac{\alpha+2}{3}} \left( \frac{v_{\text{ej}}}{0.1 \text{c}} \right)^{\frac{\alpha+2}{3}} \left( \frac{10^{-2} M_\odot}{m_{\text{ej}}} \right)^{\frac{2}{3}}. \hspace{1cm} (20)$$

These simple estimates also illustrate the dramatic effect that the changed opacities have: instead of a transient that peaks in the blue part of the optical spectrum after $\sim 1$ day (for iron-type opacities), one now expects a peak after $\sim 1$ week in the near-IR.

Early optical searches for macronova emission were carried out for several short GRBs \cite{182,192} but none of them provided convincing evidence for a macronova. The first serious evidence came from a search in the near-IR band in the aftermath of GRB 130603B using a combination of ground-based observations at $< 2$ days and HST observations at about 9 and 30 days after the burst \cite{193,194}. The difference between the observations revealed the presence of a near-IR point source at 9 days.

\footnote{See, e.g., Piran et al. (2013).}
after the GRB, this source had disappeared by the time of the second observation. These observation lead both groups\textsuperscript{183,194} to conclude that the most natural explanation was a radioactively powered transient due to decaying, freshly produced r-process elements, in other words by a macronova.

While overall compatible with the expectations for a macronova caused by the extremely neutron-rich ejecta, the inferred values\textsuperscript{183,184,194} for the ejecta masses ($m_{ej} \approx 0.03 - 0.08 \text{ M}_\odot$) and velocities\textsuperscript{195} ($v_{ej} \approx 0.2c$) are in a plausible, though slightly uncomfortably large regime for a near-equal mass nsns merger which is considered the most likely event. A natural way out would be a CBM with a mass ratio substantially different from unity, either nsns or nsbh, which would explain both larger ejecta masses and velocities\textsuperscript{112,118} Such systems, however, are expected to occur less frequently than $q \approx 1$ nsns-mergers. While r-process powered macronova are a natural interpretation, there is still room for alternative models\textsuperscript{196,197} and further macronova detections are eagerly awaited to settle the case.

5. Compact binary mergers as central engines of short gamma-ray bursts

Gamma-ray bursts are intense flashes of soft gamma-rays that are detectable approximately once per day and that reach Earth from random directions. There were early hints\textsuperscript{198–200} on a structured duration distribution, but this became only firmly established through the work of Kouveliotou et al\textsuperscript{201} who found a minimum in the duration distribution around $\sim 2$ s. They also realized that bursts below this time scale are consistently harder in their spectra than the bursts above it. This lead to the establishment of two classes of bursts, "short, hard bursts" (sGRBs; $\tau \sim 0.3$ s) and "long, soft bursts" (lGRBs; $\tau \sim 30$ s). Their spectra in the conventional $\nu - \nu F_\nu$ coordinates are usually fitted by a broken power law with a pronounced peak around $\sim 400$ KeV energies (for sGRBs; $\sim 200$ keV for lGRBs).\textsuperscript{204–206}

Compact binary mergers consisting of two neutron stars or a neutron star and a black hole have been suggested as one of the first progenitor models,\textsuperscript{109,207–211} actually long before the bimodality of the GRB distribution\textsuperscript{201} their cosmological origin,\textsuperscript{187,212,213} (and therefore the distance and energy scales) and their types of host galaxies\textsuperscript{214} were firmly established. The CBM model has survived the confrontation with three decades of new observations and many of the arguments that were brought forward in the very early papers are still considered valid today.

The discussion here will be centered around the role of compact binary mergers, for more exhaustive, general discussions of GRBs and their properties, we refer to the excellent reviews\textsuperscript{217,224} that exist on the topic.

\textsuperscript{4}This classification has, however, been subject to criticism\textsuperscript{202,203}
5.1. Confronting the CBM model with observations

We will here summarise to which extent CBMs explain the observed properties of sGRBs and where tensions between the model and observations exist. Alternative suggestions are briefly discussed in Sec. 5.5.

Energy requirements

Although often paraphrased as "the biggest explosions in the Universe" the energy requirements of sGRBs are large, but still moderate compared to the rest mass energy of a solar mass, $1 \, M_\odot c^2 \approx 1.8 \times 10^{54}$ erg. The observed, "isotropic" $\gamma$-energies (i.e. assuming that the radiation is emitted isotropically) are in the range of $E_{\gamma,\text{iso}} = 10^{48} \ldots 10^{52}$ erg\textsuperscript{225,227} which correspond to "true" energies, if the radiation is emitted into a solid angle with half-opening angle $\theta$, of $E_\gamma \approx 10^{46} \ldots 10^{50}$ erg ($\theta/8^{\circ})^2$. In fact, for the small subset of sGRBs where opening angle information is available the corrected values for both $\gamma$-ray and kinetic energies indicate $E_\gamma \sim E_{\text{kin}} \sim 10^{49}$ erg\textsuperscript{228}. The main energy reservoir that can be tapped in a CBM is the released gravitational energy which is $E_{\text{grav}} \sim GM_{\text{tot}}^2/R$, where $M_{\text{tot}}$ its the total binary mass and $R$ a separation not much larger than the Schwarzschild radius $R_s = 2GM_{\text{tot}}/c^2 = 9$ km ($M_{\text{tot}}/3M_\odot$). So that, in principle, an energy approaching $\sim M_{\text{tot}} c^2/2$, corresponding to several times $10^{53}$ erg, is available. Therefore, the energy requirements are not a serious challenge, even if only moderate efficiencies should be applicable for transforming released energy into $\gamma$-rays.

Variability time scale

The energy output of sGRBs can vary substantially on a time scale of a few ms\textsuperscript{216,227}. This is the time scale that is naturally expected from a compact binary system. Both from the dynamical time scale of a neutron star, see Eq. 4, and from the orbital period at the innermost, stable circular orbit (ISCO) of a stellar-mass black hole, see Eq. 5, one would naturally expect variations of order milliseconds.

Duration

It is commonly expected that an accretion disk is needed to transform the released gravitational binding energy into ultra-relativistic outflow and finally into the observed radiation. If one assumes that such extremely rapidly accreting systems can be described by the simple thick-disk estimate Eq. 6, there is good agreement with the typical duration of sGRBs\textsuperscript{215}, $\tau_{\text{GRB}} \sim 0.3$ s.

Host galaxies/environment

The host galaxies of short bursts are systematically different from those of long bursts. While the latter occur in unusually bright star-forming regions\textsuperscript{229,230}, the former occur in both early- ($\sim 20\%$) and late-type ($\sim 80\%$) galaxies\textsuperscript{215,223}. This points to a generally older stellar population and a substantial spread of ages. This is actually expected for compact binary mergers and has in fact been a prediction.
As discussed in Sec. 2, the inspiral time due to the emission of gravitational waves is a sensitive function of the initial orbital period $P$ and the eccentricity $e$, see Eq. (2) and Fig. 1. Since the initial values of both quantities are set by evolutionary processes of the stellar binary progenitor, one expects indeed a broad spread of inspiral times. Due to kicks that the neutron stars receive at birth large initial eccentricities may be common. If the binary as a whole receives systemic kick, substantial distances can be travelled during the inspiral time and one expects a broad distribution of merger offsets with respect to their host galaxies. The distributions obtained from binary evolution calculations agree remarkably well with those observed for sGRBs. The observed projected offsets range from 0.5 to 75 kpc with a median value of $\approx 5$ kpc. These results are interpreted as projected kick velocities from $v_{\text{kick}} \sim 20$ to $140$ km s$^{-1}$ (a median value of $\sim 60$ km s$^{-1}$) and consistent with the binary population synthesis results ($v_{\text{kick}} \sim 5$ to $500$ km s$^{-1}$) for nsns mergers. Also, the sGRB redshifts expected from CBM models are in agreement with those observed ($\sim 0.1 < z < 1.5$).

**Beaming and event rates**

The angular structure of the outflow can be inferred from so-called achromatic "jet-breaks" the light curve decay begins to steepen roughly simultaneously in all wavelength bands. This is caused by a combination of "relativistic beaming" and, in addition, by the sideways expansion of the jetted outflow. If the outflow of sGRBs is collimated into a half-opening angle $\Theta$, the emission from each emitting patch is "beamed" into a forward cone with half-opening angle $1/\gamma$, where $\gamma$ is the local Lorentz factor. As a consequence, only observers that happen to be in the beam cone can observe the emission and for them it looks like a spherical flow with the properties of the local patch. While the jet spreads sideways and the outflow becomes decelerated by the interaction with the ambient medium, the relativistic beaming angle increases, and once similar to the true geometric collimation angle, $1/\gamma \sim \Theta$, the jetted nature of the outflow becomes noticeable and a jet break in the lightcurve indicates that so far only a small patch had been visible.

If redshift, kinetic energy and ambient matter density are known, one can infer the geometric jet opening $\theta$ from the time of the jet break. Obviously, the value of this

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1Note, however, that some systems may have formed via different mechanisms. In the case of the double pulsar system, PSR J0737-3039A/B, the low space velocity, the comparatively low mass of pulsar B, the low orbital eccentricity ($e = 0.09$) and the location only $\sim 50$ pc from the Galactic midplane are all consistent with pulsar B having been formed by a mechanism different from the usually assumed core-collapse of a helium star. In fact, progenitor masses possibly substantially below $2M_\odot$ are kinematically favored.

2We refer to "collimated" to describe the geometry of the outflow and to "beamed" for the special-relativistic aberration effect.

3See, for example, Eq. (33) in Nakar (2007).
angle has severe implications for both the overall energy budget, see above, and the true event rate: for small opening angles, most bursts are unobservable and the true rate is larger by a “beaming factor” \( f_b = 4\pi/\Delta \Omega \approx 2/\theta^2 \), where \( \Delta \Omega \) is the solid angle of the beam.

So far, it has been difficult to observe jet breaks for sGRBs, mainly because of their weak afterglows. There is only a handful of cases for jet breaks and they have values of \( \sim 10^5 \) or \( 215,216 \) in rough agreement with the theoretically expected values. \( 125,247,248 \) This would imply that only roughly one out of 70 sGRBs is detectable and it would translate \( 215 \) into a true sGRB rate of \( \approx 20 \) yr\(^{-1} \) within 200 Mpc, roughly the distance accessible to Advanced LIGO/VIRGO.

"Macronova" emission
"Macronovae", see Sec. 4.4, have actually been a prediction of the compact binary merger model. The detection of the first the macronova event in the aftermath of GRB130603B \( 193,194 \) is broadly consistent with the expectations for radioactively powered transients that result from the extremely neutron-rich dynamic ejecta that produce very heavy nuclei (\( A > 130 \)). If the current ideas about other mass loss channels, see Sec. 4.1, are correct, there should be additional transients with different properties, also powered by radioactivity.

Lorentz factors
GRBs, both long and short ones, involve highly relativistic bulk flows. This was suggested as a solution of the so-called “compactness problem” \( 249,250 \) and later also confirmed by direct observations. \( 251 \) The clue to the compactness problem is to understand, how a luminous source can vary on time scales of order milliseconds and still emit optically thin radiation. Assume a (slowly moving) sphere of radius \( R \) that emits radiation and changes substantially on a time scale \( \delta t \). If it would “switch off” immediately, it would take at least the light travel time difference between closest and the farthest visible point of the sphere, \( R/c \), to convey this information to an observer. Turned around, (assuming a non-relativistic source) an observer would conclude that the source must be of a size \( R < c \delta t \approx 300 \) km(\( \delta t/1 \) ms). Now inserting typical numbers of GRBs, such a size would imply enormous optical depths (\( \tau \sim 10^{13} \)) for photons with respect to the at photon energies of \( \sim \) MeV copiously present electron-positron pairs, \( 250 \) in stark contrast to the observed, optically thin radiation. \( ^{\text{a}} \)

If the source is instead moving relativistically to the observer, the source can be larger by a factor of \( \sim \gamma^2 \) (because the observer can, due to relativistic beaming, only see a small patch of the source, but the real source size is much larger) and also the photon energy in the emission frame is lower than the observed one. The requirement \( \tau < 1 \) can then be used to place lower limits on the Lorentz factor, estimates typically favour \( 10^2 < \gamma < 10^3 \), see e.g. Lithwick & Sari \( 252 \) for a detailed

\( ^{\text{a}} \)This argument had originally been used to place an upper distance limit on GRB sources.
Multi-messenger picture of compact binary mergers

A second line of argument comes from the onset of the afterglow. Once the blast wave has deposited most of its energy in an ambient medium, it assumes a self-similar profile\(^{253}\) and X-ray and optical emissions decay as powerlaws. During this powerlaw phase the evolution of the Lorentz factor (only weakly depending on the kinetic energy of the blast and the ambient medium density) is known, and the time of the onset of this powerlaw behaviour provides a lower limit on the Lorentz factor. The presence of such large Lorentz-factors, is actually a very non-trivial constraint on GRB models. To accelerate to an asymptotic Lorentz factor of \(\gamma\), a blast of energy \(E\), cannot contain more than

\[
m_{\text{bar}} = \frac{E}{\gamma c^2} = 5.5 \times 10^{-9} M_\odot \left( \frac{E_{48}}{\gamma_{100}} \right),
\]

where we refer to a quantity \(X_n\) as \(X/10^n\) in cgs-units. It is at least not obvious, how a small fraction of the mass can receive such a disproportionate share of the released energy. How are the deposition of mass and energy so cleanly separated?

"Late-time activity"

There are also a number events that show “late-time activity” on a time scale much longer than what is naively expected from the compact binary merger model, see Sec. 2. To date this activity is still only incompletely understood.

A subset of \(\sim 20\%\) of sGRBs\(^{216,254}\) shows after the first, “prompt” spike somewhat softer, extended \(\gamma\)-ray emission that lasts for \(\sim 10 - 100\) s and is sometimes delayed in its onset. This emission component was first seen in a stacking analysis of short BATSE bursts\(^{255}\) and subsequently observed in a number of other bursts.\(^{190,256–259}\) In some cases, the fluence in this “extended emission” can exceed the one in the prompt spike, in the extreme example of GRB080503 the extended emission continues up to \(\sim 200\) s and dominates the fluence of the initial spike by a factor of 32.\(^{190}\)

Another type of late-time activity are X-ray flares following the prompt \(\gamma\)-ray emission which have actually been observed for both long an short bursts.\(^{260–266}\) In some cases, such flares have been seen many hours after the burst, e.g. GRB050724 showed a significant X-ray flare at \(\sim 14\) hours after the burst, for GRB 130603B excess X-ray emission has been detected more than a day after the main burst. For sGRBs such flares occur for different types of GRB host galaxies and both for cases with extended emission and without\(^{266}\)

Such time scales that exceed both the dynamical and the viscous time scales by many orders of magnitude are clearly uncomfortable for the compact binary merger model and had not been expected prior to detection. Given that main power source of GRBs is accretion it is natural to start from the hypothesis that it may also power this late-time activity. The dissipation parameter \(\alpha\) in Eq.\(^{6}\) is not well known, but it seems unlikely to have a value that maintains accretion for much longer than seconds. Compact binary mergers (in particular those with mass ratios deviating from unity) also eject a few percent of a solar mass into nearly unbound orbits,
and this material while providing an energy reservoir > $10^{50}$ erg provides a time scale, $\tau_{fb}$, that can easily exceed both the dynamical and the viscous time scale by orders of magnitude. However, at least in its simplest form, it may be challenged to release large amounts of energy at very late times ($\gg$ seconds). Dynamical collisions between compact objects, see Sec. 5.4, carry some promise in that respect, since each of the several close encounters produces a tidal tail that serves as a mass reservoir far from the engine. Most likely, however, they are too rare to explain the common late-time activity phenomena. While the initial, purely ballistic fallback ideas may have been too simple to explain the observations, it has become clearer over the last decade that compact binary mergers eject mass via several channels, see Sec. 4, some driven by neutrino- and/or nuclear processes, and the interaction of these different mass loss channels may yield a much more complicated picture than just the ballistic fallback. In that sense, there is some promise that also the late-time activity can be reconciled with compact binary mergers, but this question is certainly far from being understood at the moment.

5.2. How to launch ultra-relativistic outflow

Models for the ultra-relativistic outflow must –apart from questions of collimation and stability– first and foremost explain how some fraction of matter can acquire such a disproportionate share of the energy. The very large energy to rest mass ratio required to produce the inferred ultra-relativistic outflow, see Eq. (21), is a very non-trivial requirement and to date it is still a matter of debate how this is achieved in GRB. One of the mechanisms suggested to achieve this is the annihilation of neutrino anti-neutrino pairs, $\bar{\nu}_i \nu_i \rightarrow e^- e^+$. The annihilation cross-section is very small, but the neutrino luminosities are huge, see Sec. 3, and if a small fraction of the neutrino energy can be deposited in a baryon free region the typical energy of a short GRB, $\sim 10^{48}$ erg, see Sec. 5.1, could plausibly be reached. Due to the dependence of the annihilation rate on a factor $\mu \equiv 1 - \cos \theta$, where $\theta$ is the collision angle between neutrino and anti-neutrino headon collisions are favoured. Therefore, spacetime curvature or thick disk geometries can enhance the annihilation rate. A necessary requirement for neutrinos from an accretion disk is that they can escape on short enough time scale so that they are not just advected into the bh. But as estimated in Sec. 3 and further detailed by numerical simulations, the neutrino escape time scales are substantially shorter than the dynamic disk time scales.

Large neutrino luminosities, however, can also lead to strong neutrino-driven winds that can make it impossible to reach the required energy to rest mass ratio. This effect is particularly pronounced for remnants that contain a central HMNS. The situation is less critical for black hole disk systems, for which numerical simu-

\footnote{In the explicit expression there are actually two such terms, one is proportional to $\mu$, the other to $\mu^2$.}
lations\textsuperscript{117,123} find much weaker neutrino-driven winds. A recent study\textsuperscript{275} came to the conclusion that the HMNS must collapse within $\sim 100\text{ ms}$, otherwise baryonic pollution precludes the emergence of ultra-relativistic outflow.

Magnetic mechanisms are another broad class of plausible mechanisms. They are, for example, required to extract the energy of a spinning black hole via the Blandford-Znajek mechanism\textsuperscript{279} or they can more directly be transformed in outflow by buoyancy instabilities with subsequent reconnection events\textsuperscript{48,211,280} or as some form of extreme pulsar\textsuperscript{126,281–283} where the rotational energy is tapped. All these processes require very strong, near-equipartition magnetic field strengths to be effective. This is, however, likely to be achieved, at least for the short relevant time scale, since neutron stars are from the beginning endowed with strong magnetic fields and the merger process offers ample possibilities to amplify them\textsuperscript{48,130,284–287}.

Of course, it is certainly a possibility that the diversity of observed bursts is related to several of these mechanisms work in concert.

5.3. Double neutron star vs neutron star black hole binaries

Traditionally, nsns and nsbh systems have been considered as “standard GRB” progenitor, often without much distinction between them, since they were thought to both lead to the most likely engine, a stellar mass bh surrounded by an accretion disk. In recent years, however, this picture has become more differentiated. On the one hand it has become clear that bh formation in the nsns case does not have to happen early on and may actually in some cases not happen at all, see the discussion in Sec. 2. The presence of a HMNS, in turn, implies strong neutrino-
driven winds (Sec. 4.1.3), which, in turn, pose a potential threat to the ability of launching ultra-relativistic outflow. The nsbh case supposedly produces a baryon-cleaner environment (mainly due to the absence of a HMNS) and it is more likely to show imprints from jet precession. But not every combination of bh masses and spins is actually able to form a substantial accretion torus around the hole. This is mainly because the separation where the neutron star starts transferring mass/is disrupted grows with a lower power of the bh mass $M_{bh}$ than the innermost stable circular orbit (ISCO). Therefore for large enough bh masses the neutron star is swallowed before being disrupted so that no torus can form.

The final answer requires 3D numerical simulations with the relevant physics, but a qualitative idea can still be gained from simple estimates. Mass transfer is expected to set in when the Roche volume becomes comparable to the volume of the neutron star. By applying Paczynski’s estimate for the Roche lobe radius and equating it with the ns radius, one finds that the onset of mass transfer (which we use here as a proxy for the tidal disruption radius) can be expected near a separation of

$$a_{MT} = 2.17 R_{ns} \left( \frac{1 + q}{q} \right)^{1/3} \approx 26 \text{ km} \left( \frac{R_{ns}}{12 \text{ km}} \right) \left( \frac{1 + q}{q} \right)^{1/3}.$$  \hspace{1cm} (22)

Since $a_{MT}$ grows, in the limit where the mass ratio $q = M_{ns}/M_{bh} \ll 1$, only proportional to $M_{bh}^{1/3}$, but the ISCO and the event horizon grow $\propto M_{bh}$, the onset of mass transfer/disruption can take place inside the ISCO for large bh masses. At the very high end of bh masses, the neutron star is swallowed as whole without being disrupted at all. A qualitative illustration (for fiducial neutron star properties, $M_{ns} = 1.4 M_\odot$ and $R_{ns} = 12 \text{ km}$) is shown in Fig. 5 left panel. Roughly, already for black holes near $M_{bh} \approx 8 M_\odot$ the mass transfer/disruption occurs near the ISCO which makes it potentially difficult to form a massive torus from ns debris. So, low mass black holes are clearly preferred as GRB engines. Numerical simulations have shown, however, that even if the disruption occurs deep inside the ISCO this does not necessarily mean that all the matter is doomed to fall straight into the hole and a torus can still possibly form.

The bh spin also plays a crucial role for the question of disk formation, since it determines the location of the ISCO, for maximally spinning Kerr bhs, for example, it is located at only $R_{ISCO} = GM_{bh}/c^2$. This, of course, has a serious impact on the dynamics and in particular on the resulting tori. This is illustrated in Fig. 6 where two nsbh binary systems are simulated (using the SPH method) that are identical (initially circular radii with $r_c = 8GM_{bh}/c^2$ of a 1.3 $M_\odot$ ns around a spatially fixed Kerr bh) up to the bh spin parameter. In the first case of a non-rotating bh (upper row, $a_{bh} = 0$) the ns is swallowed completely within $\sim 13$ ms, while for the rapidly spinning case ($a_{bh} = 0.9$) a massive torus forms.

Fully, relativistic simulations are computationally very expensive, but it is possible to devise simple analytical models for the mass remaining outside the bh that

\footnote{This refers not only to disk mass, but also contains in addition dynamic ejecta and tidal tail.}
have a physically intuitive functional form and whose parameters can be fitted to the results of relativistic nsbh simulations. Once fitted, the results depend on the black hole mass $M_{\text{bh}}$, its dimensionless spin parameter $a_{\text{BH}}$ and the neutron star compactness $C_{\text{ns}}$. The result for the probably most likely neutron star compactness is shown in Fig. 5 right panel (their quantity $\chi_{\text{BH}}$ is what we refer to as $a_{\text{bh}}$). The results for more extreme compactness values can be found in the original paper.\cite{288} According to this model, a 10 $M_{\odot}$ bh needs to have at least a spin of $a_{\text{bh}} \approx 0.9$ to form a disk in the disruption of a “standard” neutron star, for larger/smaller stars (13.5 km/9.5 km) spin values of 0.84/0.98 are needed.

When discussing disk formation in a GRB context, it is worth keeping in mind that even seemingly small disk masses allow, at least in principle, for the extraction of energies,

$$E_{\text{extr}} \sim 1.8 \times 10^{51} \text{ erg } \left( \frac{\epsilon}{0.1} \right) \left( \frac{M_{\text{disk}}}{0.01 M_{\odot}} \right),$$

(23)

Fig. 6. Impact of black hole spin on disk formation. The panels shown in the upper row refer to the merger of a neutron star with a Schwarzschild bh ($a_{\text{bh}} = 0$), while the lower row shows the merger with a rapidly spinning Kerr bh ($a_{\text{bh}} = 0.9$). In both cases a 1.3 $M_{\odot}$ neutron star is set on a circular orbit with $r_c = 8GM_{\text{bh}}/c^2$. Due to the ISCO lying at a larger radius, for the $a_{\text{bh}} = 0$ case the black is swallowed almost completely (only very few SPH particles are left at $t = 13$ ms), while for the highly spinning case a massive accretion disk forms (last panel, second row). Note that the panels in the two rows show different scales.
that are large enough to accommodate the isotropic gamma-ray energies, \(E_{\gamma,\text{iso}} \sim 10^{50} \text{ erg}\), that have been inferred for short bursts.\(^{216}\) Collimation into half-opening angle \(\Theta\) further reduces the requirement by \(\Theta^2/2\).

If the black hole mass distribution in neutron star-black hole binary systems is indeed peaked\(^{89,90}\) around \(8 - 10 \, M_\odot\) then the large spins required to form a disk \((a_{\text{bh}} \geq 0.8)\) may seriously constrain the parameter space. Recent spin measurements in X-ray binaries\(^{250}\) find in four out of ten cases spins > 0.8, thus giving some hope that such high spins may possibly be realized also in neutron star-black hole binaries. A recent study\(^{250}\) finds that for most of their Monte Carlo models between 10 and 30\% of the neutron star-black hole mergers could be able to launch a short gamma-ray burst, but the fraction could essentially vanish for bottom-heavy black hole spin distributions. One should therefore be prepared that the contribution of neutron star-black hole mergers to the observed short gamma-ray burst rate could actually be small.

### 5.4. Gravitational wave-driven binary mergers vs. dynamic collisions

Traditionally, mostly gravitational wave-driven binary systems such as the Hulse-Taylor pulsar\(^2\)\(^,\)\(^{292}\) have been discussed as GRB engines, dynamical collisions among neutron stars and black holes in stellar systems with large number densities such as Globular Clusters have been considered a very unlikely event.\(^{293}\) More recently, however, such collisions have received a fair amount of attention.\(^{87,269,270,294–300}\) Unfortunately, their rates are at least as difficult to estimate as those of gravitational wave-driven mergers.

Collisions differ from gravitational wave-driven mergers in a number of ways. For example, since gravitational wave emission of eccentric binaries efficiently removes angular momentum in comparison to energy, primordial binaries will have radiated away their eccentricity and will finally merge from a nearly circular orbit. On the contrary, binaries that have formed dynamically, say in a globular cluster, start from a small orbital separation, but with large eccentricities and may not have had the time to circularize until merger. This leads to pronouncedly different gravitational wave signatures, “chirping” signals of increasing frequency and amplitude for mergers and initially well-separated, repeated gravitational wave bursts that continue from minutes to days in the case of collisions. Moreover, compact binaries are strongly gravitationally bound at the onset of the dynamical merger phase while collisions have total orbital energies close to zero and need to shed energy and angular momentum via gravitational wave emission and/or through mass shedding episodes in order to form a single remnant. Due to the strong dependence on the impact parameter and the lack of strong constraints on it, one therefore expects a much larger variety of dynamical behavior for collisions than for mergers.

A study that discussed various aspects of such collisions in detail\(^{296}\) came to the conclusion that such encounters could possibly produce an interesting contribution to the observed gamma-ray burst rate. Indeed, such collisions typically yield large accretion torus masses, peak neutrino luminosities\(^{87}\) in excess of \(10^{53} \text{ erg/s}\), but in some
cases substantially more, and ample possibilities to amplify initial seed magnetic fields. Perhaps most interestingly, since they typically suffer several close encounters before they can merge into a single object, and since each close encounter results in tidal tail, they produce a more complicated environment of the central merger remnant. This stores in particular substantial amounts of mass far from the remnant which only falls back on longer time scales and is maybe related to the late-time activity of some sGRBs, see Sec. 5.1. Collisions, however, also share a possible caveat with the merger cases: at least in the nsns case they likely produce very strong neutrino-driven winds that could prevent triggering a short GRB. More recently, it was concluded that collisions cannot be very frequent, since they eject very large amounts of neutron-rich matter. The dominant nsbh collision channel[7] ejects typically 0.15 $M_\odot$ per event[8], which—from cosmic nucleosynthesis arguments—restricts the (nsns + nsbh) collision rate to a maximum of $\approx 10\%$ of the estimated nsns merger rate[9]. The rare nsbh collision cases, however, would result in particularly bright “macronova” emission with

$$L_{\text{peak}} \approx 10^{41}\text{erg/s} \left( \frac{v_{\text{ej}}}{0.2c} \right)^{\frac{2}{3}} \left( \frac{10\text{ cm}^2/\text{g}}{\kappa} \right)^{\frac{2}{3}} \left( \frac{m_{\text{ej}}}{0.15M_\odot} \right)^{1-\frac{2}{3}},$$

peaking after approximately 13 days, see Eqs. (17) and (19).

5.5. Alternative models

As an alternative to the common bh-plus-disk systems, magnetars have been suggested as central engines of sGRBs, both for long and for short bursts[5,19,21,28,30,32,33,34,35]. Magnetar-like objects could be formed under a variety of circumstances. Giant flares from “young”, extragalactic magnetars with long recurrence times could plausibly be responsible for a fraction of the events observed as sGRBs[36,37,38]. A progenitor population of young magnetars, however, would be coincident with star forming regions, therefore they can be excluded as the source of the majority of sGRBs. “Old” magnetar-like objects could also be formed through the accretion induced collapse (AIC) of an accreting WD or, maybe, as a result of a double white dwarf merger[39,40]. For such old magnetars, however, no strong kicks are expected and therefore their distribution with respect to their host galaxy would differ from what is observed for sGRBs. It had also been suspected[126,283,309] that nsns mergers may in some cases result in a rapidly spinning, high-field magnetar-like object and the recent limits on the maximum mass of neutron stars $> 2M_\odot[53,54$ only make this a more intriguing possibility, see Sec. 2. A magnetar engine would have a number of obvious advantages. There would be natural formation mechanisms and a long-lived or stable central object could plausibly be responsible for the observed late-time activity. A magnetar formed as a result of a nsns-merger

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Lee et al[295] find that nsbh collisions dominate by approximately a factor of five over nsns collisions.
would obviously inherit all the benefits of the nsns merger engine model. In addition, having a magnetar at the engine of long and short bursts would naturally explain similarities between both types of bursts and with an energy emission rate (in the framework of the magnetic dipole model) of

\[ \dot{E}_{\text{md}} \propto B^2 R^6 \omega^4, \tag{25} \]

with \( B, R \) and \( \omega \) being magnetic field strength, radius and angular frequency, such an engine could naturally produce the large variety of sGRBs that is observed.

One could argue, though, that such a model, that depends on twelve powers of poorly known quantities has not much predictive power. Maybe more severe, there would be some challenge to explain both the prompt emission and the late-time activity, since the presence of a magnetar would initially produce an enormous baryon-loading which could prevent a burst to form in the first place. In fact, a recent study finds that collapse must occur before \( \sim 100 \text{ ms} \), otherwise a sGRB would be prevented. Maybe the choking of a GRB can be circumvented if the burst is only launched after a neutrino cooling time of several seconds when the neutrino-driven wind has ceased.

Another alternative engine has been suggested as a response to the discovery of late-time X-ray flares for short GRBs. In this model, the idea is that a neutron star accretes from a non-degenerate companion until it collapses into a bh surrounded by an accretion disk, generally considered the standard engine of GRBs. The X-ray flares would result when part of the relativistic ejecta interact with the extended companion star. In such a model, the neutron star would need to accrete \( \sim 0.7 M_\odot \), not an entirely trivial task given that the Eddington accretion rate is \( \dot{M}_{\text{Edd}} \sim 10^{-8} M_\odot/\text{yr} \). Another question that would need further exploration is whether/for which EOS and rotation rate combination of the collapsing neutron star a substantial disk can form outside the ISCO of the forming black hole.

In summary, there is likely room for several sGRB engines and maybe this would explain in part their diversity. However, while none of the suggestions is completely free of open questions, it is probably a fair statement that the best model to date for the bulk of sGRBs are compact binary mergers with –among them– a slight preference for nsns mergers, see the discussion in Sec. 5.3.

6. Summary

The last years have seen a tremendous progress in our understanding of compact binary mergers, both on the theoretical and the observational side, and we have provided here an overview over various of their facets. The last decade has in particular witnessed the first detection of short GRB afterglows and subsequent observations have provided a wealth of information about the environments in which short GRBs occur. Most of the observed properties find a natural explanation in the compact binary merger model, but some properties such as very late activity keep posing a problem and are not understood. They could
either point to the merger event being way more complicated than imagined in today’s models or, alternatively, to engines that are different from the standard black hole plus disk picture.

By now, practically all theoretical models agree that compact binary mergers eject enough material to be at least a major source of the heavy \((A > 130)\) r-process and nucleosynthesis calculations show good agreement with the observed abundance distributions in this regime. It has also become clear that a merger has several ways to enrich its host galaxy with neutron-rich matter: in addition to “dynamic ejecta” there are also neutrino- and/or magnetically driven winds and accretion tori that become unbound on viscous times scales and each of these channels has likely different properties. There are even recent indications that the combination of the different channels may actually not just produce the heaviest elements, but even the whole r-process range, although contributions from other sources such as a supernovae are plausible, in particular for the lighter r-process elements. A question where no consensus exists yet, is whether compact binary mergers as dominant r-process source would be consistent with the chemical evolution of galaxies. While earlier work excluded them as dominant sources, more recent studies based on hydrodynamic simulations come to the conclusion that compact binary mergers may well be consistent with the elemental scatter that is observed in stars of different ages.

Maybe one of the most exiting new developments is the recent discovery of a credible radioactively powered, electromagnetic transient candidate, also known as ”macronova” or ”kilonova”, that has been observed in the aftermath of a short GRB. In particular its time scale of about one week and the peak in the nIR are consistent with having been produced by very heavy (and therefore opaque) r-process material. If this is the correct interpretation, it connects for the first time directly short GRBs with compact binary mergers and r-process nucleosynthesis. Such transients will also increase the science returns in the era of gravitational wave astronomy that hopefully soon will begin.

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