Constraints on the merging binary neutron star mass distribution and equation of state based on the fraction of jets

Om Sharan Salafia1,2, Alberto Colombo3,2, Francesco Gabrielli4 and Ilya Mandel5,6,7

1 INAF – Osservatorio Astronomico di Brera, via E. Bianchi 46, I-23807 Merate (LC), Italy
2 INFN – Sezione di Milano-Bicocca, piazza della Scienza 2, I-20126 Milano (MI), Italy
3 Università degli Studi di Milano-Bicocca, piazza dell’Ateneo Nuovo 1, I-20126 Milano (MI), Italy
4 Scuola Internazionale Superiore di Studi Avanzati (SISSA), via Bonomea 265, I-34136 Trieste (TS), Italy
5 Monash Centre for Astrophysics, School of Physics and Astronomy, Monash University, Clayton, Victoria 3800, Australia
6 ARC Center of Excellence for Gravitational Wave Discovery – OzGrav
7 Institute of Gravitational Wave Astronomy and School of Physics and Astronomy, University of Birmingham, Birmingham, B15 2TT, United Kingdom

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ABSTRACT

A relativistic jet has been produced in the single well-localised binary neutron star (BNS) merger detected to date in gravitational waves (GWs), and the local rates of BNS mergers and short gamma-ray bursts are of the same order of magnitude. This suggests that jet formation is not a rare outcome for BNS mergers, and we show that this intuition can be turned into a quantitative constraint: at least about 1/3 of GW-detected BNS mergers, and at least about 1/5 of all BNS mergers, should produce a successful jet (90% credible level). Whether a jet is launched depends on the properties of the merger remnant and of the surrounding accretion disc, which in turn are a function of the progenitor binary masses and equation of state (EoS). The fraction of jets in the population therefore carries information about the binary distribution and EoS. Under the assumption that a jet can only be produced by a black hole remnant surrounded by a non-negligible accretion disc, we show how the jet fraction can be used to place a joint constraint on the space of BNS component mass distributions and EoS. The result points to a broad mass distribution, with particularly strong support for masses in the 1.3 – 1.6 M⊙ range. The constraints on the EoS are shallow, but we show how they will tighten as the knowledge on the jet fraction and mass distribution improve. We also discuss how to extend the method to include future BNS events, with possibly uncertain jet associations.

1. Introduction

Binary neutron star (BNS) mergers have long been considered (Eichler et al. 1989) as promising candidate progenitors of at least some gamma-ray bursts (GRBs). Evidence accumulated through the years pointing to the two observational GRB classes (‘long’ and ‘short’, Kouveliotou et al. 1993) being physically linked to two different progenitors, with collapsars (Woosley 1993) being the progenitors of long GRBs and BNS mergers (and possibly also neutron star – black hole mergers, Mochkovitch et al. 1993) being those of short GRBs (SGRBs, Berger 2014; D’Avanzo 2015). The confirmation of the collapsar progenitor scenario for long GRBs came with the association of GRB980425 with supernova SN1998bw (Galama et al. 1998). The clear evidence (Abbott et al. 2017; Mooley et al. 2018; Ghirlanda et al. 2019) in support of the presence of a relativistic jet, observed under a relatively large viewing angle with respect to its axis, in association with the GW170817 BNS merger detected by the Advanced Laser Interferometer Gravitational-wave Observatory (aLIGO, Aasi et al. 2015) and Virgo (Acernese et al. 2014) network and the fact that the inferred characteristics of such jet indicate (e.g. Salafia et al. 2019) that an on-axis observer would have observed an emission consistent with previously known SGRBs, provided the long-awaited smoking gun of the BNS-SGRB progenitor scenario. Still, the question whether all BNS mergers produce a jet (and whether other progenitors contribute a significant fraction of SGRBs) remains open. The answer to this question en-codes information about the conditions that lead to the launch of a jet and to its successful propagation up to the location where the observable emission is produced.

The conditions for the launch of a relativistic jet in the immediate aftermath of a BNS merger are set by the jet-launching mechanism, which is not entirely settled: while the Blandford & Znajek (1977) mechanism – which involves the extraction of energy from a spinning black hole by magnetic fields supported by an accretion disc – seems the most favoured one, several authors (e.g. Dai & Lu 1998; Zhang & Mészáros 2001; Metzger et al. 2008; Dall’Osso et al. 2011; Bernardini et al. 2012a; Rowlinson et al. 2013; Gompertz et al. 2013; Stratta et al. 2018; Strang & Melatos 2019; Sarin et al. 2020) invoked a rapidly spinning, highly magnetized neutron star (a proto-magnetar) to explain some observational features such as the so called ‘plateaux’ (Nousek et al. 2006) in early X-ray afterglows and the detection of ‘extended emission’ (Norris & Bonnell 2006) in gamma-rays beyond the usual short-duration prompt emission. On the other hand, many models of the early X-ray afterglow and extended emission do not require a proto-magnetar (e.g. Rees & Mészáros 1998; Zhang et al. 2006; Genet et al. 2007; Yamazaki 2009; Leventis et al. 2014; Duffell & MacFadyen 2015; Beniamini et al. 2020; Oganesyan et al. 2020; Ascenzi et al. 2020; Barkov et al. 2021; Duque et al. 2021), and the production of a successful jet by such a compact object, while possible (Usov 1992; Thompson 1994; Thompson et al. 2004; Metzger
et al. 2011; Möst et al. 2020), is theoretically disfavoured (e.g. Ciolfi 2020, see also §3.1).

Whether a BNS merger satisfies the jet-launching conditions, whatever they are, depends on the distribution of the properties of the binaries prior to merger and on the equation of state (EoS) of neutron star matter (Fryer et al. 2015; Piro et al. 2017). Despite their complexity, the merger dynamics are mostly deterministic, so the properties of the post-merger remnant and any accretion disc are directly linked to the progenitor binary masses and, in principle, spins and magnetic fields. Except for extreme cases, however, pre-merger neutron star spins and magnetic fields are thought to have limited impact (e.g. East et al. 2019; Dudi et al. 2021; Papenfort et al. 2022 for spins; e.g. Giacomazzo et al. 2009; Lira et al. 2022 for magnetic fields). Moreover, the magnetic field in the merger remnant and accretion disc is thought to undergo amplification due to Kelvin-Helmholtz instabilities and dynamo processes (e.g. Kiuchi et al. 2014, 2015; Palenzuela et al. 2021), leading to the loss of memory about the initial magnetization. Therefore, the statistical distribution of post-merger properties, and in particular the fraction of jet-launching systems in the population, is largely determined by the distribution of the progenitor binary masses and by the properties of neutron star matter, that is, by the (not yet well-known) EoS of matter beyond the nuclear saturation density. The incidence of jets in the BNS merger population therefore carries information about both binary stellar evolution and nuclear physics, which can be investigated especially through multi-messenger observations of these sources.

In this paper, we derive the posterior probability distributions on the fraction of BNS mergers that launch a jet based on the presence of a jet in GW170817 (§2.1) and on the comparison of local rate densities of SGRBs and BNS mergers (§2.2). We then describe a framework to model the jet fraction under the physically-motivated assumption that launching a jet in the aftermath of the merger requires the collapse of the remnant to a black hole on a short time scale and the presence of a non-negligible accretion disc (§3). Within this framework, we show that knowledge of the jet fraction can be used to constrain the BNS mass distribution and the EoS (§4.1). We show (§4.2) that, within this framework, the currently available information on the jet fraction leads to interesting constraints on the mass distribution, while it does not lead to informative constraints on the EoS. Nevertheless, in §4.3 we show illustrative examples of EoS constraints that can be placed in the future, when either the jet fraction or the mass distribution (or both) will be known with better precision. In Appendix A we show how the methodology can be extended to include multiple events with possibly uncertain jet associations, and in §5 we discuss our results and suggest how this methodology can be modified to become part of Bayesian hierarchical population studies of BNS mergers.

2. Jet fraction from observations

2.1. GW-detectable sub-population: binomial argument

Let us call \( f_j \) the fraction of binary neutron star (BNS) mergers in a population that produce a successful\(^1\) relativistic jet.

At the time of writing, the only binary neutron star merger detected in GWs whose localisation was sufficiently tight and close-by as to permit a thorough electromagnetic follow-up was GW170817, and a relativistic jet has been clearly detected (Abbott et al. 2017; Mooley et al. 2018; Ghirlanda et al. 2019). Let us model the production of a jet in the GW-detectable BNS subpopulation as a binomial process with single-event success probability \( f_{j,\text{GW}} \). The probability that a particular set of \( k \) events out of a total of \( n \) mergers produces a successful jet is then given by

\[
P(k \mid n, f_{j,\text{GW}}) = f_{j,\text{GW}}^k (1 - f_{j,\text{GW}})^{n-k},
\]

where the usual binomial coefficient is not present because the events can be distinguished. For a single successful event after a single observation\(^2\), this is clearly \( P(1 \mid 1, f_{j,\text{GW}}) = f_{j,\text{GW}} \).

By Bayes’ theorem, the posterior probability density of \( f_{j,\text{GW}} \) is

\[
P(f_{j,\text{GW}} \mid 1, 1) \propto f_{j,\text{GW}}^f \pi(f_{j,\text{GW}})
\]

where \( \pi(f_{j,\text{GW}}) \) is the prior probability density on \( f_{j,\text{GW}} \). An intuitive and widely adopted choice for an uninformative prior probability is a uniform distribution \( \pi(f_{j,\text{GW}}) = \pi_u = 1 \) over the domain of the parameter, which is \( f_{j,\text{GW}} \in [0, 1] \) in our case. A possibly more desirable choice is the reparametrisation-invariant ‘Jeffreys’ prior, which in this particular case reads \( \pi(f_{j,\text{GW}}) = \pi_j(f_{j,\text{GW}}) = f_{j,\text{GW}}^{1} (1 - f_{j,\text{GW}})^{-1/2} \). With the former choice, the cumulative posterior probability of \( f_{j,\text{GW}} \) is

\[
C_U(f_{j,\text{GW}} \mid 1, 1) = f_{j,\text{GW}}^f
\]

while for the latter choice it takes a slightly more complicated form, namely

\[
C_j(f \mid 1, 1) = \frac{2}{\pi} \left[ \arcsin \left( \sqrt{f} \right) - \sqrt{f - f^2} \right],
\]

where \( f = f_{j,\text{GW}} \) for short. From the cumulative posterior probability, a lower limit on \( f_{j,\text{GW}} \) at confidence level \( \alpha \) can be derived by solving \( C_j(f_{j,\text{GW}} \mid 1, 1) = 1 - \alpha \) for \( f_{j,\text{GW}} \). In the uniform prior case, this can be done analytically, yielding \( f_{j,\text{GW}} \geq \sqrt{1 - \alpha} \); for the Jeffreys prior the lower prior limit can be obtained numerically, or graphically from Figure 1, which shows a plot of \( C_U \) and \( C_j \), with the implied 90% confidence lower limit annotated. From this simple argument one can conclude, in agreement with\(^3\) Beniamini et al. (2019) and Ghirlanda et al. (2019), that the relativistic jet in GW170817 implies that a large fraction – at least about one third at the 90% confidence level – of GW-detectable BNS mergers should produce the same outcome, unless we have been extremely lucky in this first case. The corresponding 3-sigma lower limits are 5.2% (3.4%) for the uniform (Jeffreys) prior, showing that a fraction \( f_{j,\text{GW}} \) lower than several percent is highly unlikely.

\(^{1}\) In principle, a relativistic jet may be launched, but not be able to propagate through the cloud of merger ejecta and break out of it, and some authors suggested (e.g. Moharana & Piran 2017) that this is a frequent outcome. On the other hand, the expected properties of BNS merger ejecta and those of SGRB jets (also in light of GW170817) suggest that the vast majority of jets are successful in breaking out (Duffell et al. 2018; Beniamini et al. 2019; Salafia et al. 2020).

\(^{2}\) The poor localisation of the GW190425 event (Abbott et al. 2020) prevented a solid determination of the presence (or absence) of a jet, and it therefore does not contribute useful information. Nevertheless, we discuss how to account for this event (and other potentially uncertain ones) in Appendix A.

\(^{3}\) These works provide support for a large \( f_j \) based on the comparison of the single-event rate of GW170817 with the observed short gamma-ray burst rate.
2.2. Whole population: SGRB vs BNS local rate

Another route to constraining \( f_j \) is that of comparing the local rate \( R_{0,\text{SGRB}} \) of short GRBs to that of BNS mergers, \( R_{0,\text{BNS}} \). In a recent manuscript (Abbott et al. 2021b) describing compact binary merger population analyses including data from the GWTC-3 catalog (Abbott et al. 2021a), the LIGO-Virgo-KAGRA (LVK) Collaboration found local BNS rates in the range \( R_{0,\text{BNS}} \in [13 - 1900] \text{Gpc}^{-3}\text{yr}^{-1} \) (union of the 90% credible intervals obtained through different analyses), the broad range being partly due to the highly uncertain mass distribution. An independent estimate, based on a larger sample of events, can be made based on known Galactic double neutron stars (Kim et al. 2003). A recent study (Grunthal et al. 2021; see also Pol et al. 2020), which includes observational insights on the beam shape and viewing geometry of the PSRJ1906+0746 pulsar (which contributes significantly to the total rate), finds a Milky Way BNS merger rate \( R_{\text{MW}} = 32_{-9}^{+19} \text{Myr}^{-1} \). Assuming a Milky-Way-Equivalent galaxy density \( \rho_{\text{MWEG}} = 0.0116 \text{Mpc}^{-3} \) in the local Universe (Abadie et al. 2010; Kopparapu et al. 2008), this translates to \( R_{0,\text{BNS}} = 370_{-250}^{+270} \text{Gpc}^{-3}\text{yr}^{-1} \) (90% credible range, blue line in Fig. 2). In what follows, we will adopt this BNS merger rate estimate as our reference.

The local SGRB rate is poorly constrained due to the inherent difficulty in disentangling the luminosity function and redshift distribution of the cosmological population, given the current low number of events with a measured redshift and the complexity of selection effects. This emerges clearly from the diversity of local rate values present in the literature (Abbott et al. 2021b; Mandel & Broekgaarden 2021; Tan & Yu 2020; Liu & Yu 2019; Sun et al. 2017; Ghirlanda et al. 2016; Wanderman & Piran 2015; Coward et al. 2012), which is exacerbated by the difficulty in determining the average beaming factor (and its likely dependence on distance) from the available data, and by the fact that it is essentially impossible to constrain the low-end of the luminosity function (roughly below \( 5 \times 10^{49} \text{erg/s} \), as found by Wanderman & Piran 2015), making the real total rate essentially unbounded. Nevertheless, a solid lower limit on \( R_{0,\text{SGRB}} \) is set by the single-event rate of GRB170817A that is obtained by conservatively assuming it to be the only event detected by GBM over the associated sensitive volume, and neglecting the beaming factor. By carefully modelling the GBM sensitive volume for this event (as described in Appendix B) we obtain \( R_{0,\text{GRB17A}} = 34_{-16}^{+38} \text{Gpc}^{-3}\text{yr}^{-1} \) (red line in Fig. 2). Given this piece of information, we model our knowledge of \( R_{0,\text{SGRB}} \) as a log-uniform distribution bounded by \( R_{0,\text{GRB17A}} \) from below, and then compute the \( f_j \) posterior distribution as the ratio distribution of the \( R_{0,\text{SGRB}} \) and \( R_{0,\text{BNS}} \) distributions. In practice, we produce \( \{f_{j,i}\}_{i=1}^N \) samples by the following procedure: (1) we extract an \( R_{0,\text{SGRB},i} \) sample from a log-uniform distribution in the range \( 10^{-1} - 10^4 \text{Gpc}^{-3}\text{yr}^{-1} \), and an \( R_{0,\text{GRB17A},i} \) sample from the red distribution in Fig. 2; (2) if \( R_{0,\text{SGRB},i} < R_{0,\text{GRB17A},i} \) we go back to (1); (3) we extract an \( R_{0,\text{BNS},i} \) sample from the blue distribution in Fig. 2; (4) if \( f_{j,i} = R_{0,\text{SGRB},i}/R_{0,\text{BNS},i} > 1 \) we reject the sample and go back to (1), otherwise we store it. The resulting cumulative \( f_j \) distribution is shown by the grey line in Figure 1, while GW170817 ad GW190425, and thus the systematic error that stems from using the Grunthal et al. (2021) rate as a proxy for the total BNS rate is most likely less than a factor of 2.

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4 Double neutron star binaries containing an observable pulsar clearly represent only a subpopulation of merging binaries. While Grunthal et al. (2021) carefully account for selection effects and for the pulsar lifetimes in going from the observed number of such systems to the total number in the Milky Way (and to their merger rate), they do not account for a possible fraction of the population that is not currently detected by radio surveys, which may form through different channels (e.g. Vigna-Gómez et al. 2021) and may account for massive systems such as GW190425 (Abbott et al. 2020). Such systems, on the other hand, are unlikely to constitute a dominant fraction of the total (as also suggested by the relative rate densities of
two approximations of the probability density are shown in Fig. 3. This result applies to all binaries (in contrast with that obtained in the previous section, which applies to the subpopulation of GW-detectable binaries). The implied 90% credible lower limit on \( f_j \) is 21% (2.4% at 3\( \sigma \)). As shown in Fig. 3, the associated probability density is well approximated by a power law

\[
P(f_j \mid R_{0,\text{GRB17A}}, R_{0,\text{BNS}}) \propto f_j^\gamma, \tag{5}
\]

with a best fit value \( \gamma = 0.46 \).

We note that our assumption that the true local SGRB rate density is comprised between the GRB170817A local rate density (neglecting the beaming correction) and the BNS local rate density implies that we are neglecting the potential contribution of neutron star - black hole mergers (which is well justified, as the intrinsic BNS and neutron star - black hole rates are comparable – Abbott et al. 2021a – and most likely only a small fraction of the latter result in the tidal disruption of the neutron star, see e.g. Broekgaarden et al. 2021), and the log-uniform prior between \( R_{0,\text{GRB17A}} \) and \( R_{0,\text{BNS}} \) effectively corresponds to assuming a log-uniform prior on the beaming factor. Other works take the alternative route of constraining the beaming factor (or the angular dependence of the jet emission properties) from the comparison of these rates (e.g. Williams et al. 2018; Biscoveanu et al. 2020; Farah et al. 2020; Hayes et al. 2020), or of estimating the future joint BNS-SGRB detection rates as a function of such factor (e.g. Chen & Holz 2013; Clark et al. 2015).

3. Modelling the jet fraction

3.1. Blandford-Znajek jet launching conditions

The gamma-ray burst jet launching mechanism is still under debate (see Salafia & Giacomazzo 2021 for a recent discussion). In the case of a binary neutron star merger progenitor, the possible mechanisms are restricted to those compatible with the characteristics of the merger remnant: depending on the component neutron star masses and equation of state, the merger can lead either to the prompt formation (i.e. happening on a dynamical timescale) of a black hole (BH) or to a proto-neutron star (Burrows & Lattimer 1986). If the mass \( M_{\text{rem}} \) of the latter is below the maximum mass \( M_{\text{TOV}} \) that can be supported against collapse in a non-rotating neutron star, it will evolve to an indefinitely stable neutron star (NS); if the mass is above \( M_{\text{TOV}} \), but below the mass \( M_{\text{max,rot}} \) that can be supported by uniform rotation at the mass-shedding limit (Goussard et al. 1997), the remnant is said to be a ‘supra-massive’ neutron star (SMNS) and it can survive until electromagnetic spin-down eventually leads to collapse to a BH; if \( M_{\text{rem}} > M_{\text{max,rot}} \), the proto-neutron star can still be supported (Goussard et al. 1998) for a short time (typically \( \lesssim 100 \text{ ms} \)) by differential rotation before collapsing, in which case it is termed a ‘hyper-massive’ neutron star (HMNS).

Jets launched by neutron stars are observed in our Galaxy (e.g. Pavan et al. 2014; van den Eijnden et al. 2018) and several authors have argued that a rapidly spinning, highly magnetized proto-neutron star merger remnant may be able to launch a relativistic jet (e.g. Usov 1992; Thompson 1994; Thompson et al. 2004; Metzger et al. 2011; Mösta et al. 2020). On the other hand, the neutrino-driven (Dessart et al. 2009; Perego et al. 2014) and magnetically-driven (Ciolfi & Kalinani 2020) winds produced by a proto-neutron star during its early evolution are likely to load the surrounding environment with too many baryons, preventing a putative jet from reaching relativistic speeds (e.g. Ciolfi 2020).

If the remnant is a BH (or it collapses to a BH in a time much shorter than the accretion time scale) and efficient magnetic field amplification is produced by Kelvin-Helmholtz instabilities triggered at merger (Obergaulinger et al. 2010; Kiuchi et al. 2015; Aguilera-Miret et al. 2020), the Blandford-Znajek jet-launching mechanism (Blandford & Znajek 1977; Komissarov 2001; Tchekhovskoy et al. 2010) can operate, possibly enhanced by energy deposition from the annihilation of neutrino-antineutrino pairs emitted by the inner accretion disc (Eichler et al. 1989 – even though this cannot be the dominant source of jet power, as found by Just et al. 2016).

Given the difficulties with the proto-neutron star central engine, in this work we assume that only a BH remnant (possibly, but not necessarily, formed after a HMNS phase) can launch a relativistic jet. Very broadly speaking, the conditions for the Blandford-Znajek mechanism to operate are the presence of a spinning black hole (BH), a magnetised accretion disc, and a low-density funnel above the BH (along its rotation axis). Given the particular accretion conditions in the BNS post-merger phase and the expected predominantly toroidal configuration of the magnetic field in the disc (Kiuchi et al. 2014; Kawamura et al. 2016), the accretion-to-jet energy conversion efficiency \( \eta = E_{\text{jet}}/(M_{\text{disc}} c^2) \) of the Blandford-Znajek mechanism in these systems (Christie et al. 2019) is likely quite low, \( \eta \sim 10^{-3} \), as found in numerical relativity simulations (Ruiz et al. 2018) and in agreement with the estimate that can be made based on GW170817 (Salafia & Giacomazzo 2021). Assuming this efficiency, in order to produce a (short) gamma-ray burst jet with an energy in the typical (Fong et al. 2015) range \( E_{\text{jet}} \sim 10^{49} – 10^{50} \text{ erg} \), a disc mass of about \( M_{\text{disc}} \sim 0.01 – 0.1 \msun \) is needed.
Based on these arguments, we assume GRB jet launching to be possible only for HMNS or prompt BH remnants\(^5\), provided that the accretion disc mass is, conservatively\(^6\), at least \(M_{\text{disc,min}} = 10^{-3} M_\odot\). The first condition can be restated as the requirement that the gravitational mass \(M_{\text{rem}}\) of the merger remnant be at least as massive as the maximum mass that can be supported by uniform rotation \(M_{\text{max,rot}} = 1.2 M_{\text{TOV}}\), where \(M_{\text{TOV}}\) is implied by the chosen neutron star matter equation of state (EoS) and the factor 1.2 is essentially EoS-independent \(^7\) (Breu & Rezzolla 2016).

### 3.2. Determination of the remnant type

In order to compute the remnant mass \(M_{\text{rem}}\), we invoke energy conservation for a merging binary with total mass \(M = M_1 + M_2\) in the form

\[
M c^2 = M_{\text{rem}} c^2 + E_{\text{GW}} + E_{\text{disc}} + E_{\text{ej}} + E_\nu,
\]

where \(E_{\text{GW}}\) is the energy radiated in gravitational waves, \(E_{\text{disc}}\) is the total energy associated to the accretion disc and \(E_{\text{ej}}\) is that associated to the ejecta (except for disc winds, which are included in \(E_{\text{disc}}\)), and \(E_\nu\) is the energy lost in neutrinos by the remnant (in the absence of a prompt BH collapse). We compute \(E_{\text{GW}}\) using numerical-relativity-based fitting formulæ from Zappa et al. (2018), which include GW mass loss all the way from the inspiral to the early post-merger phase, and are implemented in the publicly available repository bns_lum (Bernuzzi et al. 2018). The disc energy in principle comprises its rest-mass, its gravitational potential energy and the energy associated to its rotation: the rotational and gravitational potential energy together amount approximately to 

\[
-G M_{\text{rem}} M_{\text{disc}}/2 R_{\text{disc}} = -(R_{\text{g,rem}}/R_{\text{disc}}) M_{\text{disc}} c_\text{\nu}^2/2 \quad \text{(where } R_{\text{g,rem}} \text{ and } R_{\text{disc}} \text{ are the remnant gravitational radius and the disc radius, respectively).}
\]

Since very generally \(R_{\text{disc}} \gg R_{\text{g,rem}}\), these terms can be safely neglected, keeping only the rest-mass contribution, that is \(E_{\text{disc}} \sim M_{\text{disc}} c^2\). We compute the disc rest-mass \(M_{\text{disc}}\) using the fitting formulæ proposed in Krüger & Foucart (2020). Similarly, since the typical neutron star merger ejecta velocities are subrelativistic, \(v_{\text{ej}} \sim 0.1 \sim 0.2 c\), we neglect their kinetic energy and we write \(E_{\text{ej}} \sim M_{\text{ej}} c^2\). The mass \(M_{\text{ej}}\) includes the dynamical ejecta, which we compute through the relevant fitting formulæ from Krüger & Foucart (2020), and potentially also winds from a long-lived neutron star remnant. The typical mass of neutrino-driven winds from a HMNS is expected (Dessart et al. 2009) to be on the order of \(10^{-4} M_\odot\) and can be safely neglected in our treatment; winds driven by magnetically-induced viscosity can be much more substantial, reaching \(\sim 10^{-2} M_\odot\) (Ciolfi & Kalinani 2020) in sufficiently long-lived cases, but this is still well within the expected uncertainty of the fitting formulæ we employ for the disc mass (Krüger & Foucart 2020), and therefore we do not include this component for simplicity. Similarly, since the peak neutrino luminosity of an HMNS remnant is expected to be below \(L_\nu < 10^{53} \text{erg/s}\) (Dessart et al. 2009), and the HMNS lifetime (i.e. the timescale over which differential rotation is damped) is \(t_{\text{HMNS}} \lesssim 0.1 \text{s}\), the mass lost in neutrinos \(M_\nu = E_\nu/c^2 \sim L_\nu t_{\text{HMNS}}/c^2\) is well below one percent of a solar mass, and thus we set \(E_\nu = 0\).

The employed fitting formulæ depend on the neutron star masses and on their compactness or dimensionless tidal deformability: if an EoS is assumed, the neutron star radii can be computed employing the mass ratio-dependent fitting formulæ from Zappa et al. (2018), which in-clude this component for simplicity. Similarly, since the dynamical ejecta, which we compute through the relevant fitting formulæ from Krüger & Foucart (2020), and potentially also winds from a long-lived neutron star remnant. The typical mass of neutrino-driven winds from a HMNS is expected (Dessart et al. 2009) to be on the order of \(10^{-4} M_\odot\) and can be safely neglected in our treatment; winds driven by magnetically-induced viscosity can be much more substantial, reaching \(\sim 10^{-2} M_\odot\) (Ciolfi & Kalinani 2020) in sufficiently long-lived cases, but this is still well within the expected uncertainty of the fitting formulæ we employ for the disc mass (Krüger & Foucart 2020), and therefore we do not include this component for simplicity. Similarly, since the peak neutrino luminosity of an HMNS remnant is expected to be below \(L_\nu < 10^{53} \text{erg/s}\) (Dessart et al. 2009), and the HMNS lifetime (i.e. the timescale over which differential rotation is damped) is \(t_{\text{HMNS}} \lesssim 0.1 \text{s}\), the mass lost in neutrinos \(M_\nu = E_\nu/c^2 \sim L_\nu t_{\text{HMNS}}/c^2\) is well below one percent of a solar mass, and thus we set \(E_\nu = 0\).

The employed fitting formulæ depend on the neutron star masses and on their compactness or dimensionless tidal deformability: if an EoS is assumed, the neutron star radii can be computed employing the mass radius relation, and the tidal deformabilities from the ‘C-Love’ universal relation (Yagi & Yunes 2017). This allows one to solve Eq. 6 for \(M_{\text{rem}} = M_{\text{rem}}(M_1, M_2, \Theta_{\text{EoS}})\), where \(\Theta_{\text{EoS}}\) represents a set of parameters that define the EoS. Setting \(M_{\text{rem}} < M_{\text{TOV}}\) as the condition for an indefinitely stable NS remnant; \(M_{\text{TOV}} \lesssim M_{\text{rem}} < 1.2 M_{\text{TOV}}\) as the condition for a SMNS remnant; \(1.2 M_{\text{TOV}} \lesssim M_{\text{rem}} < M_{\text{thresh}}\) as the condition for a HMNS remnant, and \(M_{\text{rem}} \sim M_{\text{thresh}}\) for direct collapse to a BH, one can then determine the fate of the merger remnant based on the initial binary masses and on the EoS. Here \(M_{\text{thresh}}\) is the threshold mass for BH direct collapse, which can be computed employing the mass ratio-dependent fitting formulæ given in Bauswein et al. 2021 (alternative recent formulæ exist, e.g. Totole et al. 2021; Kashyap et al. 2021; Kölsch et al. 2021), but we note that our results do not depend on this choice: the \(M_{\text{thresh}}\) line shown in Fig. 5 is the separation between the HMNS and BH remnant regions is only illustrative and it does not enter our assumed jet-launching conditions).

Figure 5 shows the resulting regions that correspond to different expected remnants of BNS mergers on the \((M_1, M_2)\) plane, assuming three different EoSs, namely APR4 (Akmal et al. 1998), SFHo and DD2 (Hempel et al. 2012), which predict maximum masses \(M_{\text{TOV}}/M_\odot = 2.19, 2.06\) and 2.42, respectively (all of which are above the largest Galactic neu-
In our framework, the currently available information on the EoS needs to be expressed in the form of priors on $R_{1.4}$ and $M_{\text{TOV}}$. For both quantities there exists a variety of recent constraints that use either electromagnetic or GW information, or try to combine both.

The very thorough and complete recent work by Miller et al. (2021) based on X-ray pulsar radius measurements from NICER and XMM-Newton combined with GW constraints found $R_{1.4} = 12.45 \pm 0.65 \, \text{km}$ (one sigma) consistently using

3.3. Approximate two-parameter EoS dependence

Our assumed jet-launching conditions require, in practice, the merger remnant mass to satisfy $M_{\text{rem}} > 1.2 M_{\text{TOV}}$ and the disc mass to satisfy $M_{\text{disc}} > M_{\text{disc, min}}$. A rough idea of where the former separation line stands can be obtained by writing a simplified remnant mass equation $M_{\text{rem}} \sim M - E_{\text{GW}}/c^2 - M_{\text{disc}}$, thus neglecting the dynamical ejecta mass (in addition to all other terms which we already neglect for the reasons explained in the previous section). This leads to a critical total mass for jet launch $M_{\text{crit}} \sim (1.2 M_{\text{TOV}} + M_{\text{disc}})/(1 - \eta_{\text{GW}})$, where $\eta_{\text{GW}} = E_{\text{GW}}/(M c^2)$. This places the critical mass in the range $M_{\text{crit}}/M_\odot \in [2.4, 3.26]$ for $2 \leq M_{\text{TOV}}/M_\odot \leq 2.5$, $M_{\text{disc}} \leq 0.1 M_\odot$ and $\eta_{\text{GW}} \leq 0.05$. Assuming equal masses, the component masses are therefore in the range $M_{1,2}/M_\odot \in [1.2, 1.63]$. For most viable EoSs, the NS radius is approximately constant within this mass range, so that one can quite safely assume it to be equal to $R_{1.4}$, i.e. that of a reference $1.4 M_\odot$ NS. Using this radius, the disc mass can be computed using the fitting formula from Krüger & Foucart (2020), while a reasonable estimate of $\eta_{\text{GW}}$ can be obtained as the absolute value of the Keplerian orbital energy of two point masses ($M_1, M_2$), at a $2R_{1.4}$ separation, divided by $M c^2$, which gives $\eta_{\text{GW}} \sim \nu GM/4R_{1.4} c^2$, where $\nu = M_1 M_2/M^2 = (2 + q + q^{-1})^{-1}$ is the symmetric mass ratio. This effectively reduces the EoS dependence of $M_{\text{crit}}$ to only two parameters, namely $\theta_{\text{EoS}} = \{R_{1.4}, M_{\text{TOV}}\}$. In Figure 5 we show the $M_{\text{crit}}$ obtained with this method (oblique blue dashed lines). To obtain also the line beyond which $M_{\text{disc}} < 10^{-3} M_\odot$ (horizontal blue dashed lines), we employed the relevant fitting formula from Krüger & Foucart (2020) keeping the NS radius fixed at $R_{1.4}$. These examples show that the approximate method yields results that differ from the more accurate ones (black dashed lines) at the percent level.

3.4. EoS priors and jet probability on the $(M_1, M_2)$ plane

In our framework, the currently available information on the EoS needs to be expressed in the form of priors on $R_{1.4}$ and $M_{\text{TOV}}$. For both quantities there exists a variety of recent constraints that use either electromagnetic or GW information, or try to combine both.

The very thorough and complete recent work by Miller et al. (2021) based on X-ray pulsar radius measurements from NICER and XMM-Newton combined with GW constraints found $R_{1.4} = 12.45 \pm 0.65 \, \text{km}$ (one sigma) consistently using

3.3. Approximate two-parameter EoS dependence

Our assumed jet-launching conditions require, in practice, the merger remnant mass to satisfy $M_{\text{rem}} > 1.2 M_{\text{TOV}}$ and the disc mass to satisfy $M_{\text{disc}} > M_{\text{disc, min}}$. A rough idea of where the former separation line stands can be obtained by writing a simplified remnant mass equation $M_{\text{rem}} \sim M - E_{\text{GW}}/c^2 - M_{\text{disc}}$, thus neglecting the dynamical ejecta mass (in addition to all other terms which we already neglect for the reasons explained in the previous section). This leads to a critical total mass for jet launch $M_{\text{crit}} \sim (1.2 M_{\text{TOV}} + M_{\text{disc}})/(1 - \eta_{\text{GW}})$, where $\eta_{\text{GW}} = E_{\text{GW}}/(M c^2)$. This places the critical mass in the range $M_{\text{crit}}/M_\odot \in [2.4, 3.26]$ for $2 \leq M_{\text{TOV}}/M_\odot \leq 2.5$, $M_{\text{disc}} \leq 0.1 M_\odot$ and $\eta_{\text{GW}} \leq 0.05$. Assuming equal masses, the component masses are therefore in the range $M_{1,2}/M_\odot \in [1.2, 1.63]$. For most viable EoSs, the NS radius is approximately constant within this mass range, so that one can quite safely assume it to be equal to $R_{1.4}$, i.e. that of a reference $1.4 M_\odot$ NS. Using this radius, the disc mass can be computed using the fitting formula from Krüger & Foucart (2020), while a reasonable estimate of $\eta_{\text{GW}}$ can be obtained as the absolute value of the Keplerian orbital energy of two point masses ($M_1, M_2$), at a $2R_{1.4}$ separation, divided by $M c^2$, which gives $\eta_{\text{GW}} \sim \nu GM/4R_{1.4} c^2$, where $\nu = M_1 M_2/M^2 = (2 + q + q^{-1})^{-1}$ is the symmetric mass ratio. This effectively reduces the EoS dependence of $M_{\text{crit}}$ to only two parameters, namely $\theta_{\text{EoS}} = \{R_{1.4}, M_{\text{TOV}}\}$. In Figure 5 we show the $M_{\text{crit}}$ obtained with this method (oblique blue dashed lines). To obtain also the line beyond which $M_{\text{disc}} < 10^{-3} M_\odot$ (horizontal blue dashed lines), we employed the relevant fitting formula from Krüger & Foucart (2020) keeping the NS radius fixed at $R_{1.4}$. These examples show that the approximate method yields results that differ from the more accurate ones (black dashed lines) at the percent level.
three independent frameworks, showing that the result is solid against systematics. We will therefore use a normal distribution with mean 12.45 km and standard deviation 0.65 km as our prior on $R_{1.4}$.

The available constraints on $M_{\text{TOV}}$ are more uncertain and model-dependent, as can be partly appreciated by looking at the posteriors on $M_{\text{TOV}}$ reported in Fig. 6, which have been obtained in recent works by three different groups using four different frameworks. On the other hand, a model-independent constraint is set by the mass of the heaviest known pulsar PSRJ0740+6620 (Cromartie et al. 2020), whose measurement has been recently refined (Fonseca et al. 2021) to $M_{\text{PSR}} = 2.08 \pm 0.07 M_\odot$ based on high-cadence data from the Green Bank Telescope (GBT) and the Canadian Hydrogen Intensity Mapping Experiment (CHIME). Moreover, an upper limit can be set by noting that our Blandford-Znajek jet-launching assumption implies that the total mass $M_{\text{GW17}}$ of GW170817 must have been strictly larger than 1.2$M_{\text{TOV}}$, which is the lower limit on the critical total mass for collapse that one obtains by assuming the entire gravitational mass of the binary to be retained by the remnant, without any loss in $GW_1$ ejecta, neutrinos or disc. Using these two limits, we can define our prior on $M_{\text{TOV}}$ as

$$
\pi(M_{\text{TOV}}) \propto \Phi_{\text{PSR}}(M_{\text{TOV}})(1 - \Phi_{\text{GW17}}(1.2M_{\text{TOV}})),
$$

where $\Phi_x$ represents the cumulative posterior distribution of $x$. Figure 6 shows the resulting prior (thick blue solid line), compared to recent multi-messenger constraints on $M_{\text{TOV}}$ from the literature. In practice, we construct the prior on $M_{\text{GW17}}$ using the publicly available posterior samples from Abbott et al. (2019) (we conservatively adopted the samples from the high-spin prior analysis, which yield a broader total mass posterior).

Using these $R_{1.4}$ and $M_{\text{TOV}}$ priors, we can compute the probability that a $M_1$, $M_2$ binary satisfies our jet launching conditions as

$$
P_j(M_1, M_2) = \int \int \Theta_j(M_{\text{TOV}}) \pi(R_{1.4}) dM_{\text{TOV}} dR_{1.4}
$$

where $\Theta_j(M_1, M_2, \theta_{\text{EoS}}) = 1$ if our jet launching conditions are met, and $\Theta_j = 0$ otherwise. More explicitly,

$$
\Theta_j = \Theta(M_{\text{rem}} - 1.2M_{\text{TOV}})\Theta(M_{\text{disc}} - M_{\text{disc,min}}),
$$

where $\Theta$ is the Heaviside step function, that is

$$
\Theta(x) = \begin{cases} 
0 & x < 0 \\
1 & x \geq 0 
\end{cases},
$$

and we stress that $M_{\text{rem}} = M_{\text{rem}}(M_1, M_2, \theta_{\text{EoS}})$ and $M_{\text{disc}} = M_{\text{disc}}(M_1, M_2, \theta_{\text{EoS}})$. The result of Eq. 8 for our choice of EoS priors is shown in Figure 7. The lighter the colour, the more likely our jet launching conditions are met, given our present knowledge about the EoS. For systems located in the dark lower left corner of the plot, we can confidently predict that their remnant will not collapse to a black hole on a short timescale, preventing the Blandford-Znajek mechanism from operating. The $f_j$ (or $f_{j,\text{GW}}$) lower limits from section 2 therefore lead qualitatively to the conclusion that we can exclude that the large majority of BNS systems are located in the darkest regions of this plot. The absence of observed BNS mergers or Galactic double neutron stars in these regions is therefore not due to selection effects, and models that predict a large fraction of low-mass systems, such as many population synthesis models (e.g. Fryer et al. 2012; Dominik et al. 2012; Vigna-Gómez et al. 2018; Mapelli & Giacobbo 2018), are therefore in clear tension with these results.

Fig. 6. Our prior on $M_{\text{TOV}}$ compared to recent model-dependent constraints in the literature. The blue thick line shows the $M_{\text{TOV}}$ prior we adopt in this work, which is constructed using the latest PSRJ0740+6620 mass measurement (Fonseca et al. 2021) (grey filled area) as a lower limit and the total mass of GW170817 divided by 1.2 (pink filled area) as an upper limit (see text for an explanation). The remaining coloured curves show the $M_{\text{TOV}}$ posteriors obtained recently by different groups (Raaijmakers et al. 2021a; Legred et al. 2021; Pang et al. 2021), demonstrating how the current constraints still depend sensitively on the adopted models and assumptions.

Fig. 7. Probability that a BNS would satisfy our jet launching conditions (Eq. 8) as a function of the component masses $M_1$ and $M_2$, given our present knowledge of the EoS (encoded in the priors defined in §3.4). The masses of Galactic BNS systems that merge within a Hubble time and the two GW-detected BNS are shown for comparison, in the same way as in Fig. 5.
4. Inference on the BNS mass distribution and the EoS based on the jet fraction

4.1. Derivation

Our jet launching conditions depend on the component masses $M_1, M_2$ and on the EoS through $M_{\text{TOV}}$ and the dependence of the disc and ejecta masses on the neutron star tidal deformabilities. If we were given the mass distribution $d^2P/dM_1dM_2$ of BNS mergers and the “true” EoS, we would therefore be able to compute $f_j$ directly from

$$\tilde{f}_j = \int \int \frac{d^2P}{dM_1dM_2} \Theta_1 \, dM_1 \, dM_2,$$

where we use the tilde (’$\tilde{}$’) to distinguish this functional (which depends on the mass distribution and on the EoS) from $f_j$. Similarly, we could compute the jet fraction in GW-detected binaries as

$$\tilde{f}_j,_{\text{GW}} = \int \int V_{\text{eff}} \frac{d^2P}{dM_1dM_2} \Theta_1 \, dM_1 \, dM_2,$$

where $V_{\text{eff}}$ represents the GW detector sensitive volume for a binary with masses $M_1$ and $M_2$, averaged over all extrinsic parameters. For the present detector network, whose BNS range is bound to $z \ll 1$ and whose SNR is dominated by the inspiral phase, a good approximation is $V_{\text{eff}} \propto M_\text{chirp}^{5/2}$, where $M_\text{chirp} = (M_1M_2)^{3/5}/(M_1 + M_2)^{1/5}$ is the chirp mass (e.g. Mandel et al. 2019). In what follows, we will use $f$ to indicate either $f_j$ or $f_{j,\text{GW}}$ (depending on which one of the constraints shown in Fig. 1 is used), and $\tilde{f}$ to indicate $\tilde{f}_j$ or $\tilde{f}_{j,\text{GW}}$ accordingly; we will also use $d$ to indicate the data that is used to constrain $f$, that is the information on the jet fraction in GW-detected BNS binaries or that on the SGRB and BNS local rate densities. Expressing the mass distribution in a parametric form with a set of parameters $\Theta_m$, and again the EoS with a set of parameters $\Theta_{\text{EoS}}$, we can formally use Bayes’ theorem to infer posteriors on these parameters from $f = f_{j,\text{GW}}$ or $f = f_j$, obtaining

$$P(\Theta_m, \Theta_{\text{EoS}} | f) \propto P(f | \Theta_m, \Theta_{\text{EoS}}) \pi(\Theta_m, \Theta_{\text{EoS}}),$$

where $\pi(\Theta_m, \Theta_{\text{EoS}})$ is the prior on $\Theta_m$ and $\Theta_{\text{EoS}}$, and the likelihood term is simply

$$P(f | \Theta_m, \Theta_{\text{EoS}}) = \delta(f - \tilde{f}(\Theta_m, \Theta_{\text{EoS}})).$$

Here $\tilde{f}$ represents either $\tilde{f}_j$ or $\tilde{f}_{j,\text{GW}}$, in accordance to $f$. Since the latter is not known exactly, but it is in turn inferred from $d$, we have

$$P(\Theta_m, \Theta_{\text{EoS}} | d) = \int_0^1 P(\Theta_m, \Theta_{\text{EoS}} | f) P(f | d) \, df,$$

which, using Eq. 14, simplifies to

$$P(\Theta_m, \Theta_{\text{EoS}} | d) \propto P\left( f = \tilde{f}(\Theta_m, \Theta_{\text{EoS}}) | d \right) \pi(\Theta_m, \Theta_{\text{EoS}}).$$

This equation expresses the joint constraint on the EoS and mass distribution parameter space that can be set by our arguments. For current constraints, using equations 2 and 5 and adopting a uniform prior on $f$, this can finally be written as

$$P(\Theta_m, \Theta_{\text{EoS}} | d) \propto \tilde{f}^\gamma(\Theta_m, \Theta_{\text{EoS}}) \pi(\Theta_m, \Theta_{\text{EoS}}).$$

where $\gamma = 1$ if $f = f_{j,\text{GW}}$ (see §2.1) or $\gamma = 0.46$ if $f = f_j$ (see §2.2). The somewhat shallower constraint on $f_j$ with respect to $f_{j,\text{GW}}$ is partly compensated by the $V_{\text{eff}}$ term in Eq. 12 but, as we show in the following section, the $f_{j,\text{GW}}$ is still the most constraining one between the two, given the currently available data.

4.2. Mass distribution constraints

The formalism derived in the previous section can be used to investigate the parameter space of a mass distribution model. Given our a priori knowledge of the EoS, encoded in the prior, we can marginalise Eq. 17 over $\Theta_{\text{EoS}}$ to formally obtain the posterior on the mass distribution parameters, namely

$$P(\Theta_m | d) \propto \int \tilde{f}^\gamma \pi(\Theta_{\text{EoS}}, \Theta_m) \, d\Theta_{\text{EoS}}.$$  

Taking a uniform prior on the mass distribution parameters, $\pi(\Theta_m) \propto 1$, this essentially reduces to averaging $\tilde{f}$ over the EoS uncertainty for a fixed choice of $\Theta_m$. Figure 8 shows the latter quantity for a simple mass distribution model where both binary component masses are sampled from the same Gaussian distribution with mean $\mu_m$ and $\sigma_m$, and where the marginalisation over the EoS parameters is performed using our simplified two-parameter EoS dependence (§3.3), adopting the priors introduced in §3.4, and limiting the NS masses to the range $M_{1,2} \in [1 \, M_\odot, M_{\text{TOV}}]$. Filled contours show the result obtained using $f = f_{j,\text{GW}}$, while empty contours show the corresponding result when adopting the $f = f_j$ alternative constraint. This shows that a narrow ($\sigma_m \lesssim 0.2 \, M_\odot$) Gaussian distribution with $\mu_m \lesssim 1.3 \, M_\odot$ or $\mu_m \gtrsim 1.6 \, M_\odot$ is strongly disfavoured due to the very low implied jet fraction, independently from the EoS, and consistently when using either of the constraints we considered (that is, the constraint based on the association of GW170817 to GRB170817A, which we indicate as $f = f_{j,\text{GW}}$, or the one based on the relative SGRB and BNS local rate densities, which we indicate as $f = f_j$). On the other hand, parameters representative of the main peak of the observed Galactic BNS population (white star in Fig. 8) are acceptable.

Given the fact that the analysis of GW-detected BNS (Abbott et al. 2021b; Landry & Read 2021), radio-detected Galactic BNS (Farrow et al. 2019) and both populations together (Galaudage et al. 2021) currently disfavour a narrow-peaked mass distribution, we also tested a different parametric form, that is, a power law mass distribution model where both components are sampled from a power law $P(m) \propto m^\alpha$ with $M_{\text{min}} \leq m \leq M_{\text{TOV}}$, where $m$ is the mass of either component. Figure 9 shows the resulting constraint on the $(\alpha, M_{\text{min}})$ plane, with the same colour coding as for the Gaussian mass distribution.

It is instructive to project both results on the actual mass distribution space. We stress that we assume the component masses to be sampled from the same distribution (and randomly paired) which we refer to as $P(m) = dP/dm$. The implied primary mass distribution is therefore

$$P(M_1) = \int \frac{dP}{dm}(M_1) \frac{dP}{dm}(M_2) \Theta(M_1 - M_2) \, dM_2,$$  

The priors on $\Theta_m$ and $\Theta_{\text{EoS}}$ are not necessarily independent from each other, as the mass distribution parameters can for example include $M_{\text{TOV}}$. This article is protected by copyright.
1.05 shows some improvement with respect to the priors and gen-

and similarly for the secondary mass distribution. We keep the same uniform priors as in Figures 8 and 9, that is, 1 ≤ \( \mu_m/M_\odot \leq 2 \), 0.01 ≤ \( \sigma_m/M_\odot \) ≤ 0.5, −20 ≤ \( \alpha \) ≤ 10 and 1 ≤ \( M_{\text{min}}/M_\odot \) ≤ 1.25, and we use the constraint on \( f = f_{j,GW} \). Figure 10 shows the resulting mass distribution constraints. The cut-offs below 1 M_\odot and above 2.2 M_\odot are determined by our priors; in between this range, the result shows some improvement with respect to the priors and generally disfavours narrow mass distributions. The clearest fea-

ture, common to the two parametrisations (as is particularly visible in the comparison panel), is the requirement of a non-

negligible probability for masses in the 1.3 − 1.6 M_\odot range, which are particularly well-positioned in terms of satisfying our jet-launching conditions, as shown in Figure 7. This is in good agreement with the position and width of the main mass peak as found by Galaudage et al. (2021). In Appendix C we show a comparison of these results to those from the recently circulated preprint on the LIGO/Virgo GWTC-3 population study by Abbott et al. (2021b).

4.3. EoS constraints

The posterior in Eq. 17 can also be marginalised over the mass distribution parameters, to obtain constraints on the EoS given the knowledge of the mass distribution. Unfortunately, the current uncertainties on the jet fraction and mass distribution are too large, and this simply returns the EoS priors (which are already quite informative).

To illustrate how this will improve in the future, we consider two examples. In the first one, we use \( f = f_{j,GW} \) with its current uncertainty, but we assume the intrinsic BNS masses to be well described by a Gaussian model with \( \mu_m = 1.33 M_\odot \) and \( \sigma_m = 0.09 M_\odot \). Figure 11 shows the resulting posterior distribution on the \((R_{1.4},M_{\text{TOV}})\) plane, together with the marginalised one-dimensional posterior distributions on \( R_{1.4} \) and \( M_{\text{TOV}} \), using \( f = f_{j,GW} \). We also show the marginalised likelihoods (that is, \( \int f_{j,GW} \, dM_{\text{TOV}} \) and \( \int f_{j,GW} \, dR_{1.4} \)), see Eq. 17, which best display how the jet fraction information enters the EoS parameter space. The shapes of the marginalised likelihoods are readily explained. The higher \( M_{\text{TOV}} \), the larger the fraction of SMNS remnants, which explains the trend of the \( M_{\text{TOV}} \) marginalised likelihood. Meanwhile, if \( R_{1.4} \) gets too low, many systems yield disc masses \( M_{\text{disc}} < 10^{-3} M_\odot \), which fall below our minimum requirement for a successful jet. On the other hand, a very high \( R_{1.4} \) implies large disc masses, leading to lighter remnants and therefore again to a large fraction of SMNSs. These latter two effects explain the shape of the \( R_{1.4} \) marginalised likelihood. The posterior maxima and 90% credible intervals are \( R_{1.4} = 12.3\pm0.9 \) km and \( M_{\text{TOV}} = 2.12\pm0.12 M_\odot \), demonstrating a slight improvement with respect to the priors (shown by black dashed lines in the figure), in particular for \( M_{\text{TOV}} \). We stress that these are not to be taken as actual constraints on these parameters, but as an illustrative example, since the observation of GW190425 already ruled out an intrinsic BNS mass distribution as narrow as the one employed here.

As a second example, we consider the opposite case: a still highly uncertain mass distribution, but a well constrained suc-

cessful jet fraction. We adopt here the power law mass distribu-

tion model with a uniform prior on \( M_{\text{min}} \) between 1.0 M_\odot and 1.25 M_\odot, and a uniform prior on \( \alpha \) between −10 and 10 (this is inspired by the GWTC-3 results, Abbott et al. 2021b), and assume the posterior probability on the jet fraction \( f_j \) to be represented by a Gaussian with mean \( \mu_f = 0.6 \) and standard deviation \( \sigma_f = 0.05 \). Eq. 16 then becomes

\[
P(\alpha, M_{\text{min}}, M_{\text{TOV}}, R_{1.4} \mid d) \propto \
\exp \left[ -\frac{1}{2} \left( \frac{\tilde{f}_j - \mu_f}{\sigma_f} \right)^2 \right] \pi(M_{\text{TOV}})\pi(R_{1.4})
\]

(20)

in the range where the \( M_{\text{min}} \) and \( \alpha \) priors are nonzero. The marginalised distribution \( P(M_{\text{TOV}}, R_{1.4} \mid d) \) is shown in Fig-
The result, which implies \( R_{1.4} = 12.7^{+1.00}_{-0.83} \) km and \( M_{\text{TOV}} = 2.10^{+0.13}_{-0.09} M_\odot \) (90\% credible range), shows that a good knowledge of \( f_j \) can improve our constraints on the EoS even when the BNS mass distribution is poorly known.

We conclude that in the near future, when the BNS mass distribution (plus the local rate and possibly the jet fraction) will be better constrained thanks to more GW observations, this methodology will be able to place interesting constraints on the EoS parameters and \( M_{\text{TOV}} \) in particular, to which the jet-launching fraction is most sensitive, as already found by Fryer et al. 2015.

5. Discussion and conclusions

The successful jet in GW170817 not only provided a spectacular confirmation of the long-held binary neutron star progenitor scenario for short gamma-ray bursts, but also clearly indicated that this is most likely a common outcome of this kind of mergers. In this work, we have shown how the fraction of BNS mergers that launch an ultra-relativistic jet can be used to place a joint constraint on the binary neutron star mass distribution and on the neutron star matter equation of state, under the assumption that the jet-launching mechanism requires the presence of an accretion disc with a non-negligible mass and
the collapse of the merger remnant to a black hole within a fraction of a second (i.e. a supra-massive or stable neutron star remnant are assumed not to yield SGRBs).

Based on the presence of a jet in GW170817 and adopting a simple binomial likelihood and a uniform prior, we derived a lower limit $f_j,_{GW} > 32\%$ (90\% credible level) on the fraction of jets in GW-detected binary neutron stars. By comparing the lower limit on the local short gamma-ray burst rate density that can be placed by the detection of GRB170817A by Fermi/GBM and the local binary neutron star merger rate density estimated by Grunthal et al. (2021) based on Galactic double neutron stars, we constrained the jet fraction in the whole population to be $f_j > 21\%$. This constraint is weaker than the finding of $f_j > 40\%$ (90\% credible level) by Sarin et al. (2022), a pre-print that appeared while this work was nearing submission; however, the quantitative discrepancy is likely due to the reliance by Sarin et al. (2022) on the inferred local SGRB rate density estimated by Coward et al. (2012) and Wanderman & Piran (2015) – which most likely suffer from systematics due to the uncertain luminosity function (especially at the low end) and redshift distribution of events (see §2.2) – and their assumption that promptly collapsing BNS merger products without a significant accretion disc could still launch a jet and power a SGRB.

Comparing these results with the predicted jet fraction in the population in a Blandford & Znajek (1977) jet-launching scenario, assuming equation of state priors inferred by the latest multi-messenger constraints, we can exclude mass distributions with an overwhelming fraction of low-mass ($\lesssim 1.3M_\odot$) neutron stars (as predicted by many population synthesis models, especially those which adopt the Fryer et al. 2012 ‘rapid’ supernova prescription). This finding qualitatively agrees with Sarin et al. 2022 (their Figure 1). Similarly, a mass distribution high by mass ($\gtrsim 1.6M_\odot$) components is strongly disfavoured. We caution that this conclusion depends critically on the assumed jet-launching conditions: if an alternative magnetar central engine scenario (though theoretically less favoured) was adopted, this conclusion would be significantly altered.

The method can be also used in principle to place constraints on the equation of state, but the current data are insufficient to make this methodology competitive. Nevertheless, we demonstrated that this will improve with future tighter constraints on the mass distribution and/or on the jet fraction.

In the near future, several new binary neutron stars are expected (Abbott et al. 2020) to be detected through gravitational waves in the upcoming O4 and O5 observing runs of the GW detector network (which now includes KAGRA, Somiya 2012). In addition to that, new binary pulsars are being discovered in radio surveys (Han et al. 2021; Pan et al. 2021; Good et al. 2021; Agazie et al. 2021) at an increasing rate thanks to the steady improvements in both technology and data analysis techniques. Last, but not least, our understanding of selection effects that shape the properties of the observed radio pulsar population is advancing (Chattopadhayay et al. 2020). Our knowledge of the binary neutron star merger rate and mass distribution is therefore bound to improve significantly in the next few years, leading to better prospects for this kind of methodology.

New gravitational-wave detections of binary neutron star mergers will also open the possibility of directly using the information on the presence of a jet in each event (see Appendix B for an extension to uncertain jet associations) in population studies. In a Bayesian hierarchical modelling framework such as that described by Mandel et al. (2019), the $\Theta$ term (Eq. 9) could be evaluated directly as part of single-event likelihoods, keeping the equation of state and mass distribution parameters as part of the hyperparameter vector. This can complement methods that exploit information on kilonovae ejecta masses (Li & Paczynski 1998; Metzger 2019) and SGRB afterglow observations in conjunction with gravitational-wave data (e.g. Hotokezaka et al. 2019; Radice & Dai 2019; Lazzati & Perna 2019; Barbieri et al. 2019; Coughlin et al. 2019; Dietrich et al. 2020; Breucli et al. 2021; Raaijmakers et al. 2021b).

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Appendix A: Inference on $f_{\text{jet}, \text{GW}}$ from multiple observations with uncertain jets

In this appendix, we describe how the posterior probability on $f_{\text{jet}, \text{GW}}$ can be derived in presence of multiple observations, including cases in which the presence of a jet is uncertain. Let us assume that $n$ events have been observed, and that we are able to estimate, for each $i$-th event, the probability $P_{\text{miss},i}$ that a jet would have been missed (i.e. no conclusive statement can be made) if present: this will be typically a model-dependent estimate based on the (viewing-angle-dependent) expected jet emission properties, combined with information on the fraction of the BNS merger GW localization volume that has been surveyed, and the depth of the available observations\footnote{We note that the available observations themselves could have been performed in an attempt to collect conclusive evidence in favour of a jet, in case preliminary observations provided promising indications of its presence. In that sense, $P_{\text{miss}}$ can evolve as the data are collected and, in favourable cases such as GW170817, it can get close to zero (i.e. any jet in association to GW170817, if present, would not have been missed, also factoring in the indications towards a relatively small viewing angle that stem from kilonova observations, see e.g. Breschi et al. 2021). This latter statement clearly depends also on the assumed distribution of possible jet luminosities and properties in the probed bands.}. Without using other information (such as that on the single-event masses), the intrinsic probability of the presence of a jet in each GW-detected event is just $f_{\text{jet}, \text{GW}}$. The likelihood of observing a jet in association to the $i$-th event is therefore $\bar{P}_{\text{true},i} = f_{\text{jet}, \text{GW}}(1 - P_{\text{miss},i})$, while that of observing no jet is $\bar{P}_{\text{true},i} = (1 - f_{\text{jet}, \text{GW}}) + f_{\text{jet}, \text{GW}}P_{\text{miss},i} = 1 - P_{\text{true},i}$, so that we can write

$$P(J_i | f_{\text{jet}, \text{GW}}, P_{\text{miss},i}) = J_i P_{\text{true},i} + (1 - J_i)(1 - P_{\text{true},i}),$$

(A.1)

where $J_i = 1$ if conclusive evidence in favour of a jet was found in association to the $i$-th event, and $J_i = 0$ otherwise. The $f_{\text{jet}, \text{GW}}$ posterior, up to a normalization factor, is then

$$P(f_{\text{jet}, \text{GW}} | J, P_{\text{miss}}) \propto \prod_{i=1}^{n} P(J_i | f_{\text{jet}, \text{GW}}, P_{\text{miss},i}),$$

(A.2)

where $J = \{J_1, J_2, ..., J_n\}$ and $P_{\text{miss}} = \{P_{\text{miss},1}, P_{\text{miss},2}, ..., P_{\text{miss},n}\}$. Let us evaluate Eq. A.2 in a concrete case, that of GW170817 and GW190425. In this setting, we have $n = 2$ and $J = (1,0)$. Given the short distance to GW170817, its precise localization, and the availability of an extensive multi-wavelength dataset, including VLBI observations, we can set $P_{\text{miss},1} \sim 0$. In the case of GW190425, on the other hand, the available observations can only tentatively exclude an on-axis jet (Hosseinzadeh et al. 2019). In particular, gamma-ray upper limits set by Fermi/GBM, INTEGRAL and Konus-Wind (the latter covering the entire GW localization region of GW190425) can be interpreted as the indication that the viewing angle of a putative jet, if present, must have been sufficiently large. Assuming therefore that any jet would have been missed beyond a limiting viewing angle $\theta_{v,\text{lim}}$ and calling $P(\theta_v)$ the viewing angle probability distribution of GW190425, a relativistic jet, if present, would have been missed with a probability

$$P_{\text{miss},2} \sim \int_{\theta_{v,\text{lim}}}^{\pi/2} P(\theta_v) \, d\theta_v.$$  

(A.3)

Using the probability distribution $P(\theta_v)$ reconstructed\footnote{In low signal-to-noise ratio events such as GW190425, the posterior is well-approximated by the universal inclination angle distribution of GW-detected binaries of Schutz (2011).} from the publicly available posterior samples (Abbott et al. 2020), and assuming $\theta_{v,\text{lim}} = 0.2$ rad (a representative value for an SGRB beaming angle, Fong et al. 2015), we obtain $P_{\text{miss},2} \simeq 0.94$. This implies $P(J_2 | f_{\text{jet}, \text{GW}}, P_{\text{miss},2}) \simeq 1 - 0.06 f_{\text{jet}, \text{GW}}$, which is the multiplicative correction to our result obtained by considering GW170817 only. The resulting posterior probability on $f_{\text{jet}, \text{GW}}$ is almost identical to the one obtained in §2, as can be appreciated from Figure A.1, which is unsurprising given the shallow available constraints on GW190425.

Given this framework, a relevant question that can be addressed is how many secure identifications of a jet, and/or how many tight limits on its presence, are needed to constrain $f_{\text{jet}, \text{GW}}$ to a desired precision. In that respect, we must keep in mind the fact that, while the presence of a jet can be assessed with relative certainty in favourable cases such as GW170817, its absence is generally much more difficult to prove, especially due to the strong relativistic beaming of radiation, which typically makes the detection of a putative jet extremely difficult beyond some limiting viewing angle. In well-localised cases, constraints from deep late-time radio observations can be used to probe a larger range of viewing angles, but for the majority of events this will not be a feasible route due to the combination of large localisation error box and low expected brightness of the late-time radio afterglow. This implies that it will be easier to constrain $f_{\text{jet}, \text{GW}}$ from below than from above. With these considerations in mind, we performed a simple experiment: we assumed a true value $f_{\text{jet}, \text{GW},\text{true}} = 0.5$ and constructed a sample of mock BNS events with jets assigned randomly with probability $f_{\text{jet}, \text{GW},\text{true}}$. In order to represent in a simple way the jet detectability, for each event we randomly extracted a distance and an inclination from the universal joint distribution expected for GW-detected binaries (Schutz 2011).
We then assumed a distance-dependent limiting viewing angle beyond which observations are unable to constrain the presence, or absence, of a jet. Here \(d_{\text{L,e}}\) is a parameter, which we set equal to the 25th percentile of the simulated distances, so that only for 25 percent of the events a jet can be excluded with 100% confidence if not present. If, for the \(i\)-th binary, \(\theta_i < \theta_{\text{c,lim}}\) and a jet was present, then we set \(J_i = 1\); in all other cases, we set \(J_i = 0\). For all events, we computed \(P_{\text{miss},i}\) from Equations A.3 and A.4, therefore assuming that no detection can be made beyond \(\theta_{\text{c,lim}}\), no matter how extensive the observations. Figure A.2 shows how the posterior probability on \(f_{\text{GW}}\) evolves as we include an increasing number \(N_{\text{events}}\) of these simulated events in the inference. The result suggests that several tens of detections are needed to be able to constrain the jet fraction to within 10-20%, and hundreds to get to a few percent constraint.

### Appendix B: Local SGRB rate estimate based on GRB170817A

To estimate the local SGRB rate based on GRB170817A, we model the Fermi Gamma-ray Burst Monitor (GBM) detection of GRB170817A-like sources (i.e. sources with exactly the same luminosity and intrinsic gamma-ray spectrum as GRB170817A) as a Poisson process with an expected event number \(\lambda = R_0V_{\text{eff}}T\) over a time \(T = 13\) yr (the current duration of the GBM survey), where \(V_{\text{eff}}\) is the GBM effective sensitive volume. The posterior probability on \(R_0\), assuming a single detection, is

\[
P(R_0 \mid 1) \propto P(1 \mid R_0)\pi(R_0)
\]

where \(P(1 \mid R_0) \propto R_0 \exp(-R_0 V_{\text{eff}} T)\) and we adopt the Jeffreys prior \(\pi(R_0) \propto R_0^{-1/2}\). We estimate the effective sensitive time-volume as

\[
V_{\text{eff}} T = \eta_{\text{GBM}} \int_0^\infty T \frac{dV}{dz} P_{\text{det}}(z) dz
\]

where \(\eta_{\text{GBM}} = 0.59\) accounts for the GBM field of view and duty cycle (Burns et al. 2016), \(dV/dz\) is the differential co-moving volume (Hogg 1999) and \(P_{\text{det}}(z)\) represents the probability for GBM to detect a GRB170817A-like source at a redshift \(z\). In order to compute this quantity, we assume for simplicity that the GBM detection probability depends only on the peak photon flux \(p_{64}\) of the GRB, measured in the 64 ms binning, in the 10-1000 keV band, so that \(P_{\text{det}}(z) = P_{\text{det}}(p_{64}(z))\). We then make the ansatz

\[
P_{\text{det}}(p_{64}, \alpha, p_{2}) = \frac{1}{2} \left\{ 1 + \tanh \left[ \alpha \ln \left( \frac{p_{64}}{p_{2}} \right) \right] \right\},
\]

which is a smooth function satisfying \(P_{\text{det}} = 1/2\) when \(p_{64} = p_{2}\) and \(P_{\text{det}} \sim 0\) when \(p_{64} \ll p_{2}\). The \(\alpha\) parameter controls the sharpness of the transition from 0 to 1. To fix \(\alpha\) and \(p_{2}\), we proceed as follows. We assume the intrinsic inverse cumulative distribution of short GRB photon fluxes to follow \(N_{\text{int}}(> p_{64}) \propto p_{64}^{-\alpha/2}\), as expected for uniformly distributed sources in Euclidean space, in absence of evolution of the luminosity function with distance (which is reasonable as the typical redshifts are well below 1). The observed distribution is then

\[
N_{\text{obs}}(> p_{64}, \alpha, p_{2}) \propto \int_{p_{64}}^\infty P_{\text{det}}(p, \alpha, p_{2}) dN_{\text{int}}(p) dp.
\]

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Appendix C: Comparison with GWTC-3

Figures C.1 and C.2 show comparisons of our mass distribution posteriors with the results from Abbott et al. (2021b) based on GWTC-3 (Abbott et al. 2021a) data. As a cautionary note, the definitions in our power law model match those of the POWER model from the GWTC-3 population study (Fig. C.1), but their PEAK model has the minimum mass cut $M_{\text{min}}$ as an additional free parameter (while ours is fixed at $1 M_\odot$). Despite being significantly shallower, our constraints are in general agreement with those from the population study.