Merger of Compact Stars in the Two-families Scenario

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Abstract

We analyze the phenomenological implications of the two-families scenario on the merger of compact stars. That scenario is based on the coexistence of both hadronic stars and strange quark stars. After discussing the classification of the possible mergers, we turn to detailed numerical simulations of the merger of two hadronic stars, i.e., “first family” stars in which delta resonances and hyperons are present, and we show results for the threshold mass of such binaries for the mass dynamically ejected and the mass of the disk surrounding the post-merger object. We compare these results with those obtained within the one-family scenario, and we conclude that relevant signatures of the two-families scenario can be suggested, in particular the possibility of a rapid collapse to a black hole for masses even smaller than the ones associated with GW170817; during the first milliseconds, oscillations of the post-merger remnant at frequencies higher than the ones obtained in the one-family scenario; a large value of the mass dynamically ejected; and a small mass of the disk for binaries of low total mass. Finally, based on a population synthesis analysis, we present estimates of the number of mergers for two hadronic stars, a hadronic star–strange quark star, and two strange quark stars. We show that for unequal-mass systems and intermediate values of the total mass, the merger of a hadronic star and a strange quark star is very likely (GW170817 has a possible interpretation into this category of mergers). On the other hand, mergers of two strange quark stars are strongly suppressed.

Key words: equation of state – gravitational waves – stars: neutron

1. Introduction

The first detection of gravitational waves from the merger of two compact stars\textsuperscript{8} in 2017 August (Abbott et al. 2017) represents a breakthrough for the astrophysics of compact objects and the physics of dense nuclear matter. The gravitational wave signal associated with the process of inspiral of two stars encodes information on the average tidal deformability \( \Lambda \) of the binary, which is strongly dependent on the stiffness of the equation of state (EoS) of dense matter. Several recent studies have exploited the limits on \( \Lambda \) to obtain constraints on the radii of compact stars (Abbott et al. 2018; Burgio et al. 2018; Fattoyev et al. 2018; Lim & Holt 2018; Most et al. 2018). The electromagnetic counterparts of GW170817, i.e., the short gamma-ray burst (sGRB) GRB170817A and the kilonova AT2017gfo, allow one to further constrain the EoS by indicating that the most probable merger’s remnant is a hypermassive compact star and requiring that the mass released and powering the kilonova is of the order of 0.05 \( M_\odot \) (Bauswein et al. 2017; Margalit & Metzger 2017; Annala et al. 2018; Radice et al. 2018b; Rezzolla et al. 2018; Ruiz et al. 2018). The main conclusion of those studies can be summarized by stating that the EoS of dense matter cannot be very stiff: EoSs leading to radii larger than about 13.5 km are basically ruled out. Moreover, the maximum mass of the nonrotating configuration should be less than approximately \( \sim 2.2M_\odot \), although still with a large error bar of the order of 0.2\( M_\odot \). This is an interesting result: indications of radii larger than approximately 13.5 km would indeed point to a stiff EoS in which only nucleons are present, while for smaller radii, nonnucleonic degrees of freedom could appear in compact stars.

The relevant degrees of freedom for the composition of a compact star can be inferred from the value of the radius of a star having a mass of about 1.5\( M_\odot \). First, if the radius is large, 13 km \( \lesssim R_{1.5} \lesssim 13.5 \) km, the density at the center is low enough that nonnucleonic degrees of freedom appear only in the most massive stars (Lonardoni et al. 2015). Second, if 11.5 km \( \lesssim R_{1.5} \lesssim 13 \) km, resonances (hyperons and/or deltas) will appear even for stars having masses of the order of \( \sim 1.5M_\odot \) (Maslov et al. 2015). Also, deconfined quarks can appear at the center of the star, which becomes a so-called hybrid star (HybS; Nandi & Char 2018). Finally, for \( R_{1.5} \lesssim 11.5 \) km, only disconnected solutions of the Tolman–Oppenheimer–Volkov equation exist (Alford et al. 2013; Burgio et al. 2018). These solutions are characterized by the presence of a mass window, inside which two different configurations are possible: one made only of hadronic matter and the other made partially or totally of deconfined quark matter. These configurations correspond either to the twin-stars scenario (Alford et al. 2013; Most et al. 2018; Paschalidis et al. 2018) or to the two-families scenario (Drago et al. 2014a, 2016b; Drago & Pagliara 2016; Wiktorowicz et al. 2017).

\textsuperscript{8} Here we use the generic-name compact stars to indicate NSs, which are composed only of nucleons and leptons; HSs, in which delta resonances and hyperons also take place; and strange QSs, which are entirely composed of strange quark matter.
The main differences between these two scenarios concern the mass–radius relation and the microphysics embedded in the EoS. In the twin-stars scenario, the stars containing quarks are both the most massive and the ones with the smallest radius, which reaches its minimum at the maximum mass configuration. On the other hand, in the two-families scenario, the stars having the smallest radius belong to the hadronic branch (hadronic stars (HSs)), and they have a mass of the order of (1.5–1.6)\(M_\odot\). Instead, the second family, which corresponds to strange quark stars (QSs), is characterized by not-too-small radii (of the order of or larger than about 12 km) and a maximum mass \(M_{\text{Qmax}}\) that can exceed \(2M_\odot\). Concerning the EoS, the twin-stars scenario assumes the existence of a strong first-order phase transition from hadronic matter to quark matter (often modeled via a Maxwell construction), whereas in the two-families scenario, one assumes the validity of the Bodmer–Witten hypothesis (Bodmer 1971; Witten 1984), which implies the existence of a global minimum of the energy per baryon at a finite density and for a finite value of strangeness.

The investigation of these scenarios, as well as their consequences for the interpretation of GW170817, has been discussed in previous papers (Burgio et al. 2018; Gomes et al. 2019; Nandi & Char 2018; Paschalidis et al. 2018). In particular, in Burgo et al. (2018), it was shown that the observational constraints on \(m_\Lambda\) would imply that the radius of the 1.4\(M_\odot\) configuration, \(R_{1.4}\), is larger than about 12 km within the one-family scenario, while when considering two families of compact stars, \(R_{1.4}\) could be significantly smaller, down to about 10 km.

An important point we need to stress here is that within the two-families scenario, QSs and HSs coexist. A way to classify compact objects is to look at their masses and/or radii: in a few cases, the observational indications allow one to establish to which of the two families a specific object should belong (Char et al. 2019). Once the object has been identified as a QS or an HS, one can test the model by checking whether all of the other properties are consistent with that classification. For example, the X-bursts associated with magnetar crustal modes that have been observed in a few cases are difficult to explain if those objects are QSs (see Watts & Reddy 2007), and therefore in our scheme they must be HSs. This prediction can be tested if the radius of those stars will be measured in the future. Concerning glitches, it is not yet clear whether QSs can produce such phenomena. A possibility is that at least part of the QS is in the color crystalline phase, which is rigid as well as superfluid. These are the two basic ingredients needed to model glitches; see Anglani et al. (2014) and Mannarelli et al. (2014).

An important criticism against the coexistence of QSs and HSs was put forward in Madsen (1988), where it was argued that if QSs exist, then all compact stars are QSs. This argument is based on the expectation of an abundant galactic strangelets pollution produced during the merger of two QSs. Such a flux of strangelets would trigger the conversion of all compact stars into QSs. Actually, this argument was criticized in Wiktorowicz et al. (2017), where it was shown that (i) the produced strangelets have a very large mass (order of \(10^{40}\) baryons), and thus the chance of capture by compact or main-sequence stars is very small; and (ii) smaller strangelets most likely evaporate into nucleons due to the dynamics of the merger and the strong reheating of matter after the merger; see discussion in Section 7 and N. Bucciantini et al. (2019, in preparation).

In this paper, we extend a previous preliminary work on the predictions of HS–HS mergers within the two-families scenario (Drago & Pagliara 2018). After an overview of the phenomenology of mergers within the two-families scenario, we present numerical simulations of the process of the merger of two HSs by using the Einstein Toolkit, an open-source, modular code for numerical relativity (Loffler et al. 2012). We consider two EoSs: a soft one, which contains hyperons and deltas and leads to HSs belonging to the first family, and, as a reference, a stiffer and purely nucleonic one. The main outcome of the simulations is a precise estimate of the threshold mass above which a prompt collapse occurs for the merger of two HSs. We also provide estimates of the mass ejected, which is an important quantity for the phenomenology of kilonovae. Also, we discuss in which region of the star quark nucleation can take place. Finally, we estimate the various possible merger processes within the two-families scenario by using the population synthesis code Startrack (Belczynski et al. 2002, 2008; Wiktorowicz et al. 2017).

2. EoSs of Hadronic and Quark Matter

The two-families scenario is based on the idea that in the hadronic EoS, delta resonances and hyperons do appear at large densities. To model the EoS, we adopt the relativistic mean field model SFHo of Steiner et al. (2013), with the inclusion of delta resonances and hyperons; see Drago et al. (2014b). Here we use the EoS SFHo–HD already investigated in Burgo et al. (2018) that corresponds to a coupling of delta resonances with the \(\sigma\) meson of \(x_{\sigma,\Delta} = 1.15\) (see Drago et al. 2014b; Burgio et al. 2018). A detailed motivation for this choice of the hadronic EoS will be presented in the Appendix.

For quark matter, we use the simple bag-like parameterizations of Alford et al. (2005) and Weissborn et al. (2011), where an effective bag constant \(B_{\text{eff}}\) has been introduced together with a parameter \(a_4\) that encodes pQCD corrections. By setting \(B_{\text{eff}}^{1/4} = 137.5\) MeV, \(a_4 = 0.75\), and the mass of the strange quark \(m_s = 100\) MeV, we obtain \(M_{\text{Qmax}}^0 \sim 2.1M_\odot\).

2.1. Transition from Hadronic to Quark Matter

To model the formation of quark matter, we adopt the scheme based on quantum nucleation developed in many papers; see, e.g., Heiselberg et al. (1993), Iida & Sato (1998), Berezhiani et al. (2003), Bombaci et al. (2004), and Niebergal et al. (2010). The basic idea is that the process of formation of the first drop of quark matter preserves the flavor composition, and it is therefore impossible to nucleate strange quark matter if hyperons or kaons are not already present in the hadronic phase. In Figure 1 we show the fraction \(Y_t\) of baryons as a function of the baryon density \(n_B\). Note that, as \(n_B\) increases, more and more hyperons are produced, making the conversion to strange quark matter more likely. Let us define the strangeness fraction \(Y_s = (n_s + 2(n_{\Xi^-} + n_{\Xi^0})) / n_B\). A possible way to define a threshold for the conversion of hadronic matter into strange quark matter is to request that \(d_s\), the average distance of strange quarks inside the confined phase, is of the order of or smaller than the average distance of nucleons in nuclear matter, \(d_s = n_0^{-1/3}\), where \(n_0\) is the nuclear matter saturation density, \(n_0 = 0.16\) fm\(^{-3}\). The reason is that strange quarks should be close enough to be able to mutually interact in order to help the nucleation of a first drop of deconfined quark matter. This way of defining a critical density is simple and
we expect nucleation of strange quark matter to take place: functions of the baryon density for beta-stable matter at zero temperature. The mass configuration and corresponding values of the speed of sound at the center of the maximum mass of an HS, stable with respect to quark nucleation becomes dominant has been estimated to be around 60 MeV).

In Figure 1 (upper panel), we show the fraction of the strange and not-strange baryons, together with the fraction of strangeness, at zero temperature. The previous condition on the minimal density of strangeness is satisfied above a baryon density of the order of (0.9–1) fm$^{-3}$, for which $Y_s \sim 0.2–0.3$ (see the two points on the dashed line in the figure) and which corresponds to about (6–7) $n_0$. Note that it is a density significantly larger than the threshold density of the formation of hyperons, which slightly exceeds 0.6 fm$^{-3}$, and from this viewpoint, our conditions agree with the analysis of, e.g., Bombaci et al. (2004, 2009), indicating that the first droplet of quark matter is nucleated at densities larger than the threshold density of hyperon formation. It is important to notice that the condition for quark nucleation at $T = 0$ is reached at densities smaller than those corresponding to the maximum mass mechanically stable, $M_{\text{max}}^{H_{\text{hyb}}}$. Following Bombaci et al. (2004), we therefore introduce a separate notation $M_{\text{max}}^{H_{\text{HS}}}$ for the maximum mass of an HS, stable with respect to quark nucleation.

A specific feature of the two-families scenario and the quark deconfinement mechanism is the possibility of also forming (hot) HybSs as a result of the partial turbulent conversion of hadronic matter into quark matter. These stellar objects live for a time of the order of 10 s, which is necessary for hadronic matter to convert completely into quark matter within the diffusive regime (Drago & Pagliara 2015), and they are characterized by a two-phase structure with a hot quark matter core and a hadronic matter layer (initially cooler) that are in mechanical equilibrium but out of thermal and chemical equilibrium. It is possible to construct a model for the EOS describing these transient HybSs by using Coll’s condition (Coll 1976) to define the density separating the low-density hadronic region from the high-density region already made of quarks; see Drago & Pagliara (2015) for details. The reason to discuss these short-living configurations here is that the post-merger remnant, for the cases of HS–HS and HS–QS mergers, can indeed form such HybSs in which the conversion is still proceeding (we label it Coll-hyb). Notice that since QSs are generated in this scenario by the conversion of HSs, they have a minimum mass ($M_{\text{min}}^{\text{Q}}$) that is fixed by the maximum mass of the HSs: the conservation of baryon number during the conversion implies that the total baryon number of the newly born QS is equal to the total baryon number of the progenitor HS. Since the QS is more bound than its progenitor HS, its gravitational mass is of the order of 0.1$M_{\odot}$ smaller than the gravitational mass of the progenitor, and its minimum value is $\sim (1.35–1.45)M_{\odot}$, depending on the value of the maximum mass of HSs; seeWiktorowicz et al. (2017). In turn, this implies that, at fixed gravitational mass, QSs are always larger than HSs; see red and black solid lines in Figure 2.

Another distinctive feature of the two-families scenario concerns the value of the (square) speed of sound, $c_s^2$, in dense matter: in the hadronic EOS, its value remains significantly below $\sim 0.4$ due to the several softening channels associated with the appearance of hyperons and deltas (see red line in the lower panel of Figure 1). For the quark matter EOS, within the model adopted here, at increasing density, it approaches from below the asymptotic limit of the noninteracting massless gas one-third (the conformal limit); see green line in the lower panel of Figure 1. Nevertheless, it is possible to fulfill the 2 $M_{\odot}$ limit by assuming that such massive stars are QSs. On the other hand, within the one-family scenario, the 2 $M_{\odot}$ limit requires values of the speed of sound that are significantly smaller than those corresponding to the maximum mass of an HS, stable with respect to quark nucleation.

Figure 1. Upper panel: particle fractions and strangeness fraction $Y_s$ as functions of the baryon density for beta-stable matter at zero temperature. The orange points indicate the range of densities and strangeness fraction for which we expect nucleation of strange quark matter to take place: $0.2 \lesssim Y_s \lesssim 0.3$. Lower panel: density dependence of the speed of sound for the SFHo EOS; its generalization, which includes delta baryons and hyperons SFHo-HD; and the quark matter EOS (labeled QM). The filled dots indicate the central densities and corresponding values of the speed of sound at the center of the maximum mass configuration.

Figure 2. Static and Keplerian mass–radius relations for HSs, QSs, and Coll-HybSs. We also show the results for NSs described by the SFHo EOS (the results have been obtained by using the RNS code; Stergioulas & Friedman 1995; Stergioulas 2003).
larger than the asymptotic limit, as was observed in Alford et al. (2015), Bedaque & Steiner (2015) and Chamel et al. (2013). For instance, by using the SFHi EoS, one can obtain a maximum mass of $2.06M_\odot$ (Steiner et al. 2013) with a corresponding value of the central density of $\sim 1.2$ fm$^{-3}$ and $c_s^2 \sim 0.8$; see the black line in the lower panel of Figure 1. Similarly, within the twin-star scenario, the quark matter EoS is assumed to have a large value of the speed of sound, $c_s^2 \gtrsim 0.8$ in Paschalidis et al. (2018), whereas Christian et al. (2018), Montañá et al. (2019), and Alford et al. (2019) assumed $c_s^2 = 1$. Since at asymptotically large densities, one should anyway recover the limit of $c_s^2 = 1/3$, one would need to clarify which are the physical mechanisms responsible for an initial increase of the speed of sound to values close to the causal limit and then its decrease to the conformal limit.

3. Static and Rotating Configurations

The fate of the merger of two compact stars, in particular the properties of the post-merger remnant and its emissions in gravitational and electromagnetic waves, is determined in general by two quantities: the total mass above which a prompt collapse to a black hole (BH) takes place, $M_{\text{threshold}}$, and the maximum mass of the supramassive configuration, $M_{\text{supra}}$. Both quantities strongly depend on the EoS. The fact that the outcome of a merger is (mainly) determined by those two quantities is true both in the one-family and in the two-families scenario, but in the latter, the situation is more complicated, because those quantities need to be determined for both families (and also for the short-living Coll-hyb configuration).

Let us first present the results for the structure of static and rotating stars when adopting the EoSs previously introduced. Several works have been dedicated to the study of the dependence of $M_{\text{threshold}}$ on the stiffness of the EoS: a very interesting result, presented in Bauswein et al. (2013a, 2016) and Bauswein & Stergioulas (2017) and based on explicit numerical simulations, is that the ratio $k = M_{\text{threshold}}/M_{\text{TOV}}$ scales linearly with the ratio $M_{\text{TOV}}/R_{\text{TOV}}$, i.e., with the compactness of the maximum mass configuration. Depending on the EoS, $k$ varies between 1.3 and 1.6. In Section 5 we will discuss how to estimate this parameter from direct numerical simulations of HS–HS mergers. Concerning $M_{\text{supra}}$, many studies have found that $M_{\text{supra}} \sim 1.2M_{\text{TOV}}$ (Lasota et al. 1996; Breu & Rezzolla 2016) for the case of stars with a crust (i.e., HSs and HyBs), whereas for QSs, $M_{\text{supra}} \sim 1.4M_{\text{TOV}}$ (Gourgoulhon et al. 1999; Stergioulas 2003).

We have computed both static and Keplerian configurations by using the RNS code of Stergioulas & Friedman (1995). Results are presented in Figure 2. For HSs, QSs, and neutron stars (NSs), we confirm the standard results previously described. In the two-families scenario, we have also computed $M_{\text{Coll-hyb}}$ for our Coll-hyb configurations, which turns out to be of the order of $2.6M_\odot$, respecting the general relation between $M_{\text{TOV}}$ and $M_{\text{supra}}$ of compact stars with a crust. As discussed above, while Coll-hyb configurations are chemically and thermally out of equilibrium, they represent a necessary intermediate stage of the evolution of the post-merger remnant in the two-families scenario. Their stability is therefore important both for HS–HS mergers and for HS–QS mergers. We have not computed the maximum mass of a differentially rotating Coll-hyb configuration $M_{\text{Coll-hyb, threshold}}$ through direct numerical simulations, but we can safely assume that $M_{\text{threshold}} \geq M_{\text{supra}}$.

For the sake of the following discussions, let us summarize the possible values of the masses discussed above. First, let us introduce the quantity $M_{\text{tot}}$, which is the sum of the gravitational masses of the two objects undergoing the merger:

$$M_{\text{tot}} = M_1 + M_2.$$  \hspace{1cm} (1)

When discussing the critical masses, one has to be careful to distinguish between $M_{\text{tot}}$ and the gravitational mass that remains after mass ejection (and gravitational wave emission) from the system, $M_{\text{remnant}}$. The relation between these two masses is of the form

$$M_{\text{remnant}} = (1 - \alpha)M_{\text{tot}}.$$  \hspace{1cm} (2)

where $\alpha$ is of the order of a few percent. Plausible values of the masses discussed above are

1. HSs:
   \begin{align*}
   M_{\text{TOV}}^H & \sim M_{\text{max}}^H \sim 1.6M_\odot, \\
   M_{\text{supra}}^H & \sim 1.9M_\odot, \\
   M_{\text{threshold}}^H & \sim 2.5M_\odot; \\
   
   
   \end{align*}

2. QSs:
   \begin{align*}
   M_{\text{TOV}}^Q & \sim 2.1M_\odot, \\
   M_{\text{supra}}^Q & \sim 3M_\odot, \\
   M_{\text{threshold}}^Q & \sim M_{\text{supra}}^Q \sim 3M_\odot; \text{10}
   
   \end{align*}

3. Coll-HyBs:
   \begin{align*}
   M_{\text{Coll-hyb}}^\text{Coll-hyb} & \sim M_{\text{TOV}}^\text{Coll-hyb} \sim 2.1M_\odot, \\
   M_{\text{Coll-hyb}}^\text{supra} & \sim 2.6M_\odot, \\
   M_{\text{Coll-hyb}}^\text{Coll-hyb, threshold} & \sim 1.5M_{\text{TOV}}^\text{Coll-hyb} \sim 3.1M_\odot.
\end{align*}

Concerning HSs, within the two-families scenario, there is an uncertainty on $M_{\text{TOV}}^H$ (whose value determines $M_{\text{supra}}^H$ and $M_{\text{threshold}}^H$) and $M_{\text{max}}^H$ due to the difficulty in estimating the critical value of the central density at which the formation of quark matter is triggered. A possible range of values for $M_{\text{max}}^H$ is $M_{\text{max}}^H = (1.5 - 1.6)M_\odot$, and $M_{\text{TOV}}^H$ takes only slightly higher values.

The value of $M_{\text{TOV}}^Q$ has a larger uncertainty stemming from the unknown value of the maximum mass of compact stars, which, as mentioned in the Introduction, is estimated to be $(2.2 \pm 0.2)M_\odot$. Interestingly, this same range is the one allowed by the microphysics of nucleation of quark matter in hadronic matter. For significantly larger values of $M_{\text{TOV}}^Q$ the process of nucleation of quark matter would no longer be energetically convenient (Drago et al. 2019).

Finally, the uncertainty on $M_{\text{TOV}}^\text{Coll-hyb}$ translates into a similar uncertainty on the values of $M_{\text{supra}}^\text{Coll-hyb}$ and $M_{\text{threshold}}^\text{Coll-hyb}$. Once a value for $M_{\text{TOV}}^\text{Coll-hyb}$ is chosen, the value of $M_{\text{supra}}^\text{Coll-hyb}$ can be explicitly computed, while the value of $M_{\text{threshold}}^\text{Coll-hyb}$ has only been estimated using the results of Weih et al. (2018), indicating that the maximum mass of differentially rotating stars is a factor of $\sim 1.5$ larger than $M_{\text{TOV}}^\text{Coll-hyb}$.

Recently, a reanalysis of the observational values of masses and radii within the two-families scenario was performed in Char et al. (2019). It was shown that all existing constraints can

\[^{10}\text{In Bauswein et al. (2009), two quark EoSs have been investigated with } M_{\text{TOV}}^Q \sim 1.64M_\odot \text{ and } 1.88M_\odot, \text{ respectively. From Figure 2 of that paper, one can extract the ratio between } M_{\text{threshold}}^\text{Coll-hyb} \text{ and } M_{\text{TOV}}, \text{ which turns out to be close to the ratio between } M_{\text{supra}}^\text{Coll-hyb} \text{ and } M_{\text{TOV}}.\]
be satisfied within the model. Interestingly, the ranges of the masses of HSs and QSs turn out to be in agreement with the distribution of the masses of the two subpopulations discussed in Antoniadis et al. (2016), Alting et al. (2018), and Tauris et al. (2017). The two distributions are centered at $m_1 = 1.39M_\odot$ with a dispersion $\sigma_1 = 0.06M_\odot$, and at $m_2 = 1.81M_\odot$ with a dispersion $\sigma_2 = 0.18M_\odot$, respectively (Antoniadis et al. 2016). Our values for $M_{\text{max}}^H$ and $M_{\text{max}}^Q$ are very close to the upper and lower limits for $m_1$ and $m_2$, respectively. It is therefore tempting to interpret the first subpopulation as composed of HSs and the second of QSs.

4. Classification of the Mergers

Within the two-families scenario, there are three possible merger events: HS–HS, HS–QS, and QS–QS. Notice that HS–HS is obviously the only possible type of merger in a one-family scenario and that QS–QS has been discussed in the past in a few papers by various authors without making any assumption on the number of families of compact stars (Haensel et al. 1991; Bauswein et al. 2009, 2010). The HS–QS merger is instead possible only within the two-families scenario. These different possibilities are displayed in Figure 3 in which the two axes correspond to the two parameters determining the main properties of the binary, i.e., the chirp mass $M_{\text{chirp}}$ and the mass asymmetry $q = m_2/m_1$ (for the sake of discussion, we have made a specific choice of the maximum and minimum masses of HSs and QSs; see the caption). It is clear how rich the phenomenology of mergers can be within the two-families scenario; for instance, within the small window labeled “All-comb.,” there could be events of HS–HS, HS–QS, and QS–QS merger. On the other hand, there are regions of the diagram in which we predict only one specific type of merger to be possible (for instance, only HS–HS mergers to the left of the green line). As a consequence, for the same value of $M_{\text{chirp}}$, there could be a prompt collapse to a BH if the merger is of the HS–HS type, while a hypermassive or supramassive configuration could form if it is an HS–QS merger; see Drago & Pagliara (2018).

Note that the classification presented in Figure 3 depends only on the maximum and minimum masses of the HSs and QSs, since it does not suggest the outcome of the mergers (this question will be discussed in Section 4.2) but only indicates which types of mergers are possible within the two-families scheme. Finally, the probabilities of the various merger processes will be addressed in Section 6 by using a population synthesis code.

4.1. Quark Deconfinement in the Merger of Compact Stars

To better understand why the properties of the hybrid configurations Hyb-Coll are also important to establish the fate of a merger in the two-families scenario, one has to recall in which way the process of quark deconfinement starts and proceeds inside a compact star. Quark deconfinement will evolve differently in the three types of merger we are discussing.

1. HS–HS merger. In this case, quark matter is not present in the system before the merger. As already discussed in the previous section, the trigger for quark matter nucleation is the presence of hyperons, since they provide a finite density of strange quarks. Once the density/temperature at the center of a compact object exceeds the critical values at which the density of the hyperons is large enough, the process of quark deconfinement starts. This implies that quark deconfinement begins immediately after the merger of two HSs, because a few milliseconds after the merger, the density and temperature increase, and hyperons start forming at the center of the star (Sekiguchi et al. 2011). In Section 5.6 we will investigate the conditions for quark nucleation by studying the results of the merger simulations. The process of quark deconfinement is initially very rapid, since it is accelerated by hydrodynamical instabilities, and in a few milliseconds, it converts the bulk of the star into quarks. After this initial phase, the instabilities are suppressed, and the process becomes much slower: the conversion of the most external layer into quarks can take about $10^{-5}$ s.11 As in Section 3.1, it is therefore very important to estimate the maximum mass of the stable hybrid hadron–quark configuration that forms a few milliseconds after the merger, $M_{\text{TOV}}^\text{hyb}$, and its rotating counterparts, $M_{\text{TOV}}^{\text{hyb, Coll}}$ and $M_{\text{TOV}}^{\text{hyb, Coll, supr}}$, because their values determine the fate of the newly formed object after the central region has deconfined.

2. HS–QS merger. In this case, quark matter is already present in the system. Notice that at the moment, we are not able to simulate the merger HS–QS because it involves two different EoSs. We can, however, compare this process with the HS–HS one by noticing that in the HS–QS case, a large fraction of the star is already in the quark matter phase. The process of rapid combustion of hadrons into quarks can still take place if the quark matter phase does not already occupy the whole region that

11 Notice that the rapid burning stops at an energy density that satisfies Coll’s condition (Drago & Pagliara 2015). Since the temperatures reached by the system after the merger and therefore the energy density distribution depend on the pre-merger configuration, the value of $M_{\text{TOV}}^{\text{hyb}}$ is not “universal” but rather depends on the pre-merger history. The value we have provided is therefore just an approximation; the real value could be established through a simulation of the merger in which the process of quark deconfinement is described as, e.g., in the numerical simulation of Pagliara et al. (2013), but this is at the moment a very challenging numerical problem.
satisfies Coll’s condition. In this case, it is also possible that the process of quark deconfinement is significantly faster, because the highly turbulent initial post-merger phase can mix quark and hadronic matter so that the mixing area is much larger than in the laminar case, and the process is thus greatly accelerated. Also in this case, \( M_{\text{coll-hyb}}^{\text{TOV}} \), \( M_{\text{supra}}^{\text{TOV}} \), and \( M_{\text{coll-hyb}}^{\text{TOV}} \) play an important role in the determination of the fate of the merger.

3. QS–QS merger. This merger has been discussed in a few papers in the past (see, for instance, Bauswein et al. 2010). The main reason for the interest in this process is the possibility of a “clean” environment, since it was assumed that little or no hadronic matter was emitted during and after the merger (Haensel et al. 1991). Although this is not strictly true (quite a significant amount of matter can be ejected dynamically at the moment of the merger), the system is more “clean” of baryon contamination, since significantly less matter can be ejected by, e.g., neutrino ablation. This is due to the much larger binding energy of strange quark matter with respect to nuclear matter. Later, we will discuss the possible phenomenological signatures of this type of merger. It is clear that since in this case, all hadronic matter is already made of deconfined quarks, the only relevant quantities are \( M_{\text{coll-hyb}}^{\text{TOV}} \), \( M_{\text{supra}}^{\text{TOV}} \), and \( M_{\text{coll-hyb}}^{\text{TOV}} \).

4.2. Possible Outcomes of a Merger

Let us discuss the possible outcomes of a merger with the associated phenomenology. A first crucial problem we need to address is to clarify in which way an sGRB can be produced in the two-families scenario.

4.2.1. sGRB Inner Engine

There are mainly two ways to generate an sGRB: one is based on the formation of a BH (MacFadyen & Woosley 1999) and the other on the formation of a protomagnetar (Metzger et al. 2011). These two mechanisms have something in common: they both need to wait for a long enough time so that the environment becomes less baryon-loaded and a jet can be launched. The ultimate source of energy to power the sGRB is the accretion disk around the BH, in the first case, and the rotational energy in the case of the protomagnetar. If the sGRB is due to the formation of the BH, the duration of the prompt emission is regulated by the duration of the disk around the BH. Since the amount of material in the disk is significantly less in the case of mergers than in the case of collapsars, one can explain why the durations of long GRBs and sGRBs differ by roughly 2 orders of magnitude. It is more problematic to explain that difference within the protomagnetar mechanism, since the strength of the magnetic field and the rotational frequencies are comparable in the two cases (Rowlinson et al. 2013). In Drago et al. (2016a), it was suggested that the duration of sGRBs is linked to the time needed to deconfine the surface of the star close to the rotation axis: as long as nucleons are still present on that surface, the baryonic load is too large to launch a jet, while when quarks occupy the whole surface, the baryonic load rapidly drops to zero. The duration of the burst is therefore related to the time needed to completely deconfine the surface of the post-merger compact object.

In our scenario, we assume that both mechanisms are possible: if a BH forms a few tens of ms after the merger (once differential rotation has been dissipated\(^{12}\)), an sGRB can be launched by using the energy extracted from the accretion disk, while if the remnant is stable for at least a few seconds (as for a supramassive star), it can generate an sGRB within the protomagnetar mechanism. In this second scenario, an extended emission could also be obtained with a duration that is connected to the time needed for the supramassive star to collapse (Rowlinson et al. 2013; Lü et al. 2015).

Let us now discuss the outcomes of the mergers in the three possible cases within the two-families scenario. For all of the critical masses, we will adopt the values discussed in Section 3.

4.2.2. HS–HS Merger

The HS–HS merger is the case examined in the numerical simulations of this paper. In the following sections, we will determine the value of \( M_{\text{tot}}^{\text{threshold}} \), and we will investigate whether the conditions for quark nucleation can be reached. Depending on the value of \( M_{\text{tot}} \), the outcomes of the mergers are the following.

\[
(1) \quad M_{\text{tot}} > M_{\text{coll-hyb}}^{\text{TOV}} \quad \text{Within the parameters’ space discussed above, such a possibility is never realized.}
\]

\[
(2) \quad M_{\text{tot}} < M_{\text{coll-hyb}}^{\text{TOV}} \quad \text{Within the parameters’ space discussed above, such a possibility is never realized.}
\]

\[
(3) \quad M_{\text{tot}} \leq \min[M_{\text{supra}}^{\text{TOV}} \quad (1 - \alpha) \quad M_{\text{coll-hyb}}^{\text{TOV}}] = M_{\text{coll-hyb}}^{\text{TOV}} \quad \text{There is a direct collapse to a BH.}
\]

\[
(4) \quad M_{\text{tot}} > M_{\text{supra}}^{\text{TOV}} \quad \text{This case is similar to the analogous case of HS–HS merger. A possible difference could come from}
\]

\(^{12}\)The time needed for the dissipation of differential rotation is rather uncertain, since it depends on the difficult estimates of the viscosity in the system. Recently, it has been suggested that the hypermassive configuration could last for a timescale of 1 s (Gill et al. 2019).

\(^{13}\)For highly asymmetric systems (\( q = 0.8 \)), even in the cases of prompt collapse, an sGRB could be released due to the formation of a massive torus around the BH, which could have a mass of a few 0.1\( M_{\odot} \); see Rezzolla et al. (2010) and Giacomazzo et al. (2013). Since the torus mass scales linearly with \( M_{\text{TOV}} \) (which is \( \sim 1.6 \) for HSs), we expect that asymmetric HS–HS systems that collapse promptly to a BH would leave tori significantly lighter than the ones produced within the one-family scenario (for which \( M_{\text{TOV}} \) must be larger than 2\( M_{\odot} \)).
the inspiral gravitational wave’s signal, since we expect a larger value of $\tilde{\Lambda}$ with respect to the case of HS–HS systems.

(2) $M_{\text{Coll-hyb}}^{\text{hyb}}/(1 - \alpha) < M_{\text{tot}} < M_{\text{Coll-hyb}}^{\text{hyb}}$. The remnant is a hypermassive star. Within the two-families scenario, this is the only possible interpretation of the event of 2017 August. We will discuss that event separately in Section 7.

(3) $M_{\text{tot}} < M_{\text{Coll-hyb}}^{\text{hyb}}/(1 - \alpha)$. The phenomenology is similar to case (3) of HS–HS merger, and it is characterized by the possibility of producing an extended emission following the sGRB.

4.2.4. QS–QS Merger

(1) $M_{\text{tot}} > M_{\text{threshold}}^Q \sim M_{\text{supra}}^Q/(1 - \alpha)$. This case is most likely irrelevant from the phenomenological point of view, since the mass distribution of binary systems peaks at $1.33M_\odot$ with $\sigma \sim 0.11M_\odot$ (Kiziltan et al. 2013) and the value of $M_{\text{supra}}^Q \sim 3M_\odot$. Anyway, its phenomenology would be very similar to those of the previously discussed cases of prompt collapse.

(2) $M_{\text{tot}} < M_{\text{threshold}}^Q \sim M_{\text{supra}}^Q/(1 - \alpha)$. This case was explored a long time ago by Haensel et al. (1991) as a possible mechanism for producing an sGRB. In particular, it was recognized that the environment around the merger is relatively clean of baryon pollution, and therefore it is easier to produce ejecta with a high Lorentz factor. The ultimate source of the sGRB in their analysis is provided by the cooling of the remnant, and the spectra would therefore be very similar to those of a blackbody. Notice, however, that rotation and magnetic field were not considered in that pioneering study. If the post-merger object rotates rapidly and develops a strong magnetic field, then the protomagnetar mechanism could also be applied to this merger, and an extended emission could also be generated, as in the other cases discussed above. A completely open issue concerns the possible presence and features of a kilonova associated with such an event. In a recent study by Paulucci et al. (2017), it was proposed that the evaporation of strangelets could be very efficient and that r-process nucleosynthesis could be at work, although not for the production of lanthanides. The kilonova associated with that event could thus show a dominant blue component, but further and more detailed studies are necessary, since the chain of nucleosynthesis could also possibly be completed in this case.

5. Numerical Simulations of the Merger of HS–HS

We considered binary NS merger processes using numerical relativity in the case of pure HSs and for two EoSs, namely, the SFHo and SFHo-HD EoS, presented in Section 2 and we studied the behavior of equal-mass models, as summarized in Table 1, for both of them.

5.1. Numerical Methods and Initial Data

The numerical methods used to perform these simulations are the same as those of De Pietri et al. (2016, 2018) and Maione et al. (2016, 2017). Here we report the general simulation setup and parameters and refer to those previous articles for more details. In particular, the resolution used in this work is $dx = 0.1875$ CU $= 277$ m (see De Pietri et al. 2016 for a discussion of the convergence properties of the code).

The simulations were performed using the Einstein Toolkit (Löffler et al. 2012), an open-source modular code for numerical relativity based on Cactus (Allen et al. 2011).

The evolved variables were discretized on a Cartesian grid with six levels of fixed mesh refinement, each using twice the resolution of its parent level. The outermost face of the grid was set at $720M_\odot$ (1040 km) from the center. We solved the BSSN-NOK formulation of Einstein’s equations (Nakamura et al. 1987; Shibata & Nakamura 1995; Baumgarte & Shapiro 1998; Alcubierre et al. 2000, 2003) implemented in the Mclachlan module (Brown et al. 2009) and the general relativistic hydrodynamics equations with high-resolution shock-capturing methods implemented in the publicly available module GRHydro (Baiotti et al. 2005; Möst et al. 2014). In particular, we used a finite-volume algorithm with the HLLE Riemann solver (Harten et al. 1983; Einfeldt 1988) and the WENO reconstruction method (Liu et al. 1994; Jiang & Shu 1996). The combination of the BSSN-NOK formalism for Einstein’s equations and the WENO reconstruction method was found in De Pietri et al. (2016) to be the best setup within the Einstein Toolkit, even at low resolutions. For the time evolution, we used the method of lines with fourth-order Runge–Kutta (Runge 1895; Kutta 1901). For numerical reasons, the system is evolved on an external-matter atmosphere set to $\rho_{\text{atm}} = 6.1 \times 10^3$ g cm$^{-3}$ (as in Lehner et al. 2016) that it is just slightly larger than the one used in Sekiguchi et al. (2015) and Radice et al. (2018), but it is still sufficiently low to avoid the atmosphere’s inertial effects on the ejected mass on a timescale of $\approx 10$ ms after the onset of the merger (Sekiguchi et al. 2015). Initial data were generated with the LORENE code (Gourgoulhon et al. 2001) as irrotational binaries in the conformal thin sandwich approximation. In this work, we analyze different combinations of equal-mass setups for binary systems using the SFHo-HD (Drago et al. 2014b; Burgio et al. 2018) and SFHo EoSs (Steiner et al. 2013). For both models, the initial distance was set to 44.3 km, like in Maione et al. (2016) and Feo et al. (2017). Matter description is performed using a 10 piece piecewise polytropic approximation for an EoS of the type $P = P(\rho, e)$ supplemented by an additional thermal component described by $\Gamma_{\text{th}} = 1.8$. The parameterization of the EoS uses the following prescriptions:

$$\epsilon = \epsilon_0(\rho) + \epsilon_{\text{th}},$$

$$p = p_0(\rho) + (\Gamma_{\text{th}} - 1)\rho\epsilon_{\text{th}},$$

where $\epsilon_{\text{th}}$ is an arbitrary function of the thermodynamical state that has the property of being zero at $T = 0$, and $\epsilon_0(\rho)$ and $p_0(\rho)$ are the internal energy and pressure at $T = 0$. In particular, for a piecewise polytropic approximation with $N$ pieces, we can write

$$p_0(\rho) = K_i\rho^{\Gamma_i},$$

$$\epsilon_0(\rho) = \epsilon_i + \frac{K_i}{\Gamma_i - 1}\rho^{\Gamma_i^{-1}},$$

and all of the coefficients are set once the polytropic index $\Gamma_i$ ($i = 0, \ldots, N - 1$), the transition density $\rho_i$ ($i = 1, \ldots, N - 1$), and $K_0$ are chosen. One should note that here $\epsilon_{\text{th}}$ plays the role of temperature, and it may be converted (using its interpretation as a perfect fluid thermal component) to a temperature scale as $T = (\Gamma_{\text{th}} - 1)m_b\epsilon_{\text{th}}$, where we assumed a conventional value of $m_b = 940$ MeV$/c^2$ for the mass of a free baryon.
5.2 Gravitational Waves Extraction

During the simulations, the gravitational wave signal is extracted using the Newman–Penrose scalar $\Psi_4$ (Newman & Penrose 1962; Baker et al. 2002) using the module WeylScalar4. The scalar is linked to the gravitational wave strain by the following relation, which is valid only at spatial infinity:

$$\Psi_4 = \hat{h}_+ - i \hat{h}_x,$$

where $h_+$ and $h_x$ are the two polarizations of the gravitational wave strain $h$. The signal is then decomposed into multipoles using spin-weighted spherical harmonics of weight $-2$ (Thorne 1980). This procedure is made by the module MULTIPLE using the following relation:

$$\psi_4(t, r, \theta, \phi) = \sum_{l=2}^{\infty} \sum_{m=-l}^{l} \psi_{4lm}(t, r, \theta, \phi) Y_{lm}(\theta, \phi).$$

In this work, we focus only on the dominant $l = m = 2$ mode; therefore, we will refer to $h_{22}$ as $h$ for the rest of this paper. In order to extract the gravitational wave strain from $\Psi_4$ and minimize the errors due to the extraction, one has to extrapolate the signal extracted within the simulation at a finite distance from the source to infinity, in order to satisfy the previous equation. Then the extrapolated $\Psi_4$ is integrated twice in time, employing an appropriate technique in order to reduce the amplitude oscillations caused by the high-frequency noise aliased in the low-frequency signal and amplified by the integration process (Reisswig & Pollney 2011; Nakano et al. 2015). The procedure adopted in this work is extensively discussed in Maione et al. (2016). First, $\Psi_4$ is extrapolated to spatial infinity using the second-order perturbative correction of Nakano et al. (2015),

$$R \psi_{4lm}(t_{\text{ret}})|_{r=\infty} = \left(1 - \frac{2M}{R}\right)(R R_{4lm}(t_{\text{ret}})) - \frac{(l - 1)(l + 2)}{2R} \hat{h}_{4lm}(t_{\text{ret}}) + \frac{(l - 1)(l + 2)(l^2 - 1)}{8R^2} \hat{h}_{4lm}(t_{\text{ret}}),$$

where the gravitational wave strains $\hat{h}_{4lm}(t_{\text{ret}})$ at finite radius (in our case, at a fixed coordinate radius $R = 1033$ km) are computed by integrating the Newman–Penrose scalar twice in time with a simple trapezoid rule, starting from zero coordinate time, and fixing only the two physically meaningful integration constants.
material leaving a coordinate sphere surface with a given radius using the \( u_t < -1 \) condition. Using this procedure, we determined the unbound matter flow for a set of different radii from 65 CU \( \simeq 96 \text{ km} \) up to 700 CU \( \simeq 1034 \text{ km} \), which is the border of the computational domain.

5.4. Results of the Numerical Simulations

In this work, we simulate equal-mass binary systems for HS–HS and NS–NS configurations during about four orbits before merger and for about 20 ms after merger. All of the simulated models show a very similar behavior. To show the main properties of the evolution, snapshots for model SFHo-HD 118vs118 are shown in Figures 4 and 5. In particular, Figure 4 shows the density of the expelled matter and its temperature at different times, while in Figure 5 its velocity and the localization of unbound mass are shown at the same times. The computed gravitational wave signal of our simulation is shown in Figures 6 and 7, where we display the power spectrum extracted using a detector at \( \mathcal{R} = 1034 \text{ km} \) for models described by the SFHo-HD EoSs and the SFHo EoS, respectively. Associated with each simulation, we report the values of \( \mathcal{M}_{\text{tot}} \) and \( \mathcal{M}_{\text{disk}} \) in Table 2. Their values are also plotted in Figure 12 as a function of \( \mathcal{M}_{\text{tot}} \). A striking feature of the gravitational wave spectrum is that models with \( \mathcal{M}_{\text{tot}} \geq 2.48 \) and \( \mathcal{M}_{\text{tot}} \geq 2.84 \) for SFHo-HD and SFHo, respectively, do not show any post-merger \( f_2 \) peak. This is related to the fact that, in these cases, we have a direct collapse to a BH (less than 1 ms after the merger). The fact that a prompt collapse occurs is clearly shown in the plot of the maximum density as a function of time for models described by the SFHo-HD (Figure 8) and SFHo (Figure 9) EoSs. In the cases in which a prompt collapse does not occur, the main structure of the spectrum is the same as the one discussed in Takami et al. (2015), Bauswein & Stergioulas (2015), and Maione et al. (2017), and references therein. Importantly, the post-merger spectrum is characterized by a main \( f_2 \) peak whose frequency depends on the EoS. Moreover, its frequency increases with the mass of the star, while its tidal parameter decreases. One should also notice that secondary peaks are present, but their relative importance decreases as the total mass of the system is reduced.

The presence of a well-defined main peak and its frequency \( f_2 \) could lead to a clear determination of the kind of merger and the properties of the EoS. In fact, that would be a precise signature that the system did not directly collapse to a BH and that its spectrum encodes information on the properties (mainly the stiffness) of dense matter. For example, the merger of two compact stars of mass 1.18 \( \mathcal{M}_\odot \) shows a main peak at frequency \( f_2 = 3.71 \) or 2.88 kHz if the EoS describing the matter is SFHo-HD or SFHo, respectively. It is interesting to remark that the value of this frequency is strictly related to the value of \( \mathcal{R}_{\text{th}} \) as discussed in Bauswein et al. (2016). Therefore, our result does not depend on the details of the EoS but on the request of having very small radii in the hadronic branch. Unfortunately, the amount of gravitational wave energy available for the mode detection (see Table 2) is going to be at best of the order of 0.1 \( \mathcal{M}_{\text{tot}} \); therefore, it is unlikely that such a peak will be detected by the present generation of gravitational wave detectors, but they will be a distinct feature to be observed in third-generation detectors.

Let us now discuss the results for the ejected mass and the amount of mass that is left in the disk. In Figure 10 we show an example of the outflow for SFHo-HD 118vs118. By following...
the flow of the unbound matter during the evolution, we can estimate the ejected mass by integrating (at each radius) on time the unbound matter flow. In this model, the total flows of unbound mass at coordinate radii of 96, 220, 443, and 590 km are (5.0, 10.7, 11.24, 11.45) $\times 10^6$ $M_\odot$, respectively. Indeed, this procedure for computing the total ejected mass is very sensitive to the extraction radius, and this is a common property of all of the simulated models. In particular, we compared the results obtained by using this method with those obtained by considering the maximum of total unbound mass present on the whole computational domain at a given time. We note that the difference between these two estimates of $M_{ej}$ is always less than about 20%, and the two results get closer by increasing the extraction radius. From the lower panel of Figure 10 one should note that most of the dynamically unbound mass is generated between 100 and 200 km from the center of the remnant, and it is related to three main outbursts of ejected matter associated with the three bouncings of the maximum density. This is more clear looking at Figure 11 where we display the radial density profile (as a function of the coordinate radius) of the unbound matter at different times. There, one should note that three shocks are created within a 100 km radius from the center of the star ($t = 0.14, 1.15$, and $2.15$ ms), the first being of lower amplitude. Then they move out, amplify, and spread while the system evolves. This is analogous to what is observed for the flux of matter at different radii (see Figure 10, upper panel), where one can see that at $R = 100$ km, one has three clearly separated peaks that join and spread as the unbound matter flows outside the computational domain.

The same overall dynamics is associated with the mass ejection for all of the considered models (except models that feature a direct collapse to a BH and for which the mass...
ejection is suppressed). All of the results for $M_{\text{ej}}$ and $M_{\text{disk}}$ as a function of $M_{\text{tot}}$ are summarized in the lower and upper panels, respectively, of Figure 12. The explicit values are also reported in Table 2. A striking feature concerning the mass ejected is the presence of a maximum, which is located at a value of $M_{\text{tot}}$ slightly smaller than the value of $M_{\text{threshold}}$. This is one of the main results of the present work. In particular, this maximum is located at $M_{\text{2.72}}$ for the SFHo EoS, and it corresponds to an ejected mass of $\approx 16 M_{\odot}$. For the SFHo-HD case, we have not determined the maximum of the ejected mass, which should be located below or close to $\approx 2.36 M_{\odot}$ (the simulation with the lowest $M_{\text{tot}}$). For this specific value, we have found an ejected mass of $\approx 13 M_{\odot}$, very close to the maximum for the SFHo EoS. The plots of the ejected mass show a very steep decay of its value as the total mass of the binary increases, and the ejection is almost completely suppressed in the case of a prompt collapse to a BH.

Another clear difference between the two EoSs concerns the value of $M_{\text{disk}}$: in the case of SFHo-HD, it is always $\lesssim 0.01 M_{\odot}$, whereas for SFHo, it can be an order of magnitude larger. These are potentially very interesting results for the phenomenology of the kilonova. As we will discuss in Section 8, the luminosity of the different components of the kilonova (red, blue, and purple kilonova; Perego et al. 2017) strongly depends on the specific mechanisms for the ejection of mass during the different stages of the merger. A relevant fraction of the ejected mass can, in principle, come from the disk; therefore, a very small value of $M_{\text{disk}}$ corresponds to a strong suppression of certain components of the kilonova signal.

Figure 5. Snapshots of the $XY$ and $XZ$ projections of the time component of the fluid velocity, $u_t$, and the matter velocity, $v$, for the SFHo-HD 118vs118 model. In each panel, the white region outside the solid cyan line refers to the atmosphere, which was excluded from the computation, while the colored lines refer to the density contours that correspond to the densities used for the polytropic approximation discussed in Section 5. The solid white lines in the first two rows refer to the $u_t = -1$ condition, which defines the difference between bound and unbound matter, as discussed in Section 5.3. The time $t$ is relative to the merger time for this model.
Let us now compare our results with the ones obtained by similar numerical simulations present in the literature. In particular, there are a few simulations using the SFHo EoS for masses similar to our model SFHo 136vs136. The outflow dynamics for this model is shown in Figure 13. The results we obtain are in general agreement with the results obtained in Sekiguchi et al. (2015, 2016) for the SFHo EoS. In those works, it was found that the associated ejected mass is of the order of $10 \, \text{m} \, M_\odot$ for the mass of the two stars of $1.35 \, M_\odot$ and of the same order of magnitude for unequal-mass systems of approximately the same total mass (Shibata et al. 2017). In the same work, it was also found that the expected disk mass (for the same systems) is of the order of $50 \text{–} 120 \, \text{m} \, M_\odot$ and indeed of the same order of magnitude as the values found in the present work. Computations of the ejected mass for the SFHo EoS were also presented in

![Figure 6](image6.png) Spectrum of the gravitational wave signal for models described by the SFHo-HD EoS. All of the given models show a reduced $f_2$ post-merger peak that is fully suppressed for models that have a direct collapse to a BH ($M_{\text{BH}} \geq 2.48 M_\odot$).

![Figure 7](image7.png) Spectrum of the gravitational wave signal for models described by the SFHo EoS. Note that the model with a total mass of $2.84 \, M_\odot$ has no post-merger peak. The model with a total mass of $2.80 \, M_\odot$ shows just a marginal $f_2$ peak, while all of the less massive models show a distinct main $f_2$ post-merger peak and two (lower-amplitude) side peaks (forming a three-peak structure) that become less important as the total mass decreases.

![Figure 8](image8.png) On the left, we show the density for which the strangeness fraction is $0.2$ (black line) and $0.3$ (red line) as a function of the temperature in MeV. On the right, we show the maximum mass density as a function of time for the models described by the SFHo-HD EoS.

| Model          | $M_{eJ}$ (m$M_\odot$) | $M_{\text{disk}}$ (m$M_\odot$) | $E_{\text{POST}}$ (m$M_\odot$) | $f_2$ (kHz) | $t_{\text{BH}}$ (ms) |
|----------------|------------------------|-------------------------------|-------------------------------|------------|----------------------|
| SFHo-HD 118vs118 | 12.993                 | 12.92                         | 25.42                         | 3.71       | 3.82                 |
| SFHo-HD 120vs120 | 9.435                  | 13.81                         | 22.42                         | 4.00       | 3.16                 |
| SFHo-HD 122vs122 | 4.290                  | 8.34                          | 6.06                          | ...        | 1.91                 |
| SFHo-HD 124vs124 | 3.011                  | 2.89                          | 0.66                          | ...        | 1.00                 |
| SFHo-HD 126vs126 | 0.737                  | 2.45                          | 0.20                          | ...        | 0.79                 |
| SFHo-HD 128vs128 | 0.055                  | 0.74                          | 0.04                          | ...        | 0.70                 |
| SFHo-HD 130vs130 | 0.043                  | 0.71                          | 0.01                          | ...        | 0.59                 |
| SFHo-HD 118vs118 | 1.968                  | 76.66                         | 42.16                         | 2.88       | ...                  |
| SFHo-HD 120vs120 | 2.085                  | 71.72                         | 43.87                         | 2.90       | ...                  |
| SFHo-HD 122vs122 | 1.730                  | 91.81                         | 42.00                         | 2.90       | ...                  |
| SFHo-HD 124vs124 | 1.824                  | 65.58                         | 52.98                         | 2.96       | ...                  |
| SFHo-HD 126vs126 | 2.375                  | 60.86                         | 58.33                         | 2.98       | ...                  |
| SFHo-HD 128vs128 | 3.145                  | 112.24                        | 50.33                         | 3.05       | ...                  |
| SFHo-HD 130vs130 | 4.523                  | 73.82                         | 59.33                         | 3.06       | ...                  |
| SFHo-HD 132vs132 | 6.007                  | 88.87                         | 67.29                         | 3.18       | 25.75                |
| SFHo-HD 134vs134 | 9.511                  | 49.27                         | 65.09                         | 3.25       | 13.55                |
| SFHo-HD 136vs136 | 16.244                 | 30.71                         | 58.76                         | 3.40       | 9.42                 |
| SFHo-HD 138vs138 | 10.367                 | 16.09                         | 46.06                         | 3.55       | 5.06                 |
| SFHo-HD 140vs140 | 4.170                  | 6.45                          | 22.39                         | ...        | 2.13                 |
| SFHo-HD 142vs142 | 2.247                  | 2.01                          | 2.02                          | ...        | 0.98                 |
Bauswein et al. (2013b) using SPH dynamics, Lehner et al. (2016), and, using the WhiskyTHC code, Bovard et al. (2017) and Radice et al. (2018a), where a comparison among all of these numerical results is shown in their Table 3. One can notice that the variability among the various estimates is rather large. Concerning our work, in particular, a few potentially relevant effects have not been incorporated, such as a treatment of the neutrino transport and a fully consistent description of the thermal component of the EoS, and we use a piecewise polytropic approximation. However, our results are in very good agreement with Sekiguchi et al. (2015), where full beta equilibrium, thermal evolution, and approximate neutrino cooling and absorption were considered.

Figure 9. Plot of the maximum density as a function of time for models described by the SFHo EoS. It should be noted that the first model showing no bouncing corresponds to a total mass of $2.84 M_\odot$.

Figure 10. Upper panel: estimates of the unbound material using the $v_r < -1$ condition from the flow crossing spherical surfaces with different coordinate radii. We show the results at four radii. We note that there are three peaks at the closest surface, which correspond to the “rebounds” of the hypermassive NS formed after the merger. Those peaks become larger when crossing the farther surfaces due to the different velocities of the ejected material. Lower panel: time-integrated fluxes at different crossing radii. The black line corresponds to the total unbound mass on the whole grid at a given time.

Figure 11. Snapshots at different times of the radial density of the unbound matter as a function of the radius for the SFHo-HD 118vs118 model. The appearance of various peaks at a given radius and for different times corresponds to the shocks generated by the oscillations of the remnant displayed in Figure 8.

Figure 12. Representation of the ejected and disk mass as reported in Table 2 in terms of the total mass of the model. The blue stars refer to the SFHo-HD models, while the red dots refer to the SFHo ones. The dotted vertical lines correspond to the values of $M_{\text{threshold}}$ for both EoSs.
Another important effect that one should consider is related to the dependence of \( M_{\text{ej}} \) and \( M_{\text{disk}} \) on the mass asymmetry \( q \). This issue has been discussed in Rezzolla et al. (2010), Giacomazzo et al. (2013), and, more recently, Kiuchi et al. (2019). It is expected that \( M_{\text{disk}} \) is always larger for unequal-mass binaries because of the more efficient tidal interactions and angular momentum transfer during the merger. The same consideration seems not to apply in general for \( M_{\text{ej}} \), and further studies are needed, in particular when considering models close to the threshold mass, because the state after the merger may rapidly change from producing a direct collapse to forming a short-lived remnant (Kiuchi et al. 2019).

5.5. Estimate of the Threshold Mass

A first estimate of the values of the threshold mass for the SFHo and SFHo-HD EoSs can be obtained by using the empirical formulæ presented in Bauswein et al. (2013a, 2016) and Bauswein & Stergioulas (2017). The ratio \( k = M_{\text{threshold}} / M_{\text{TOV}} \) scales linearly with the compactness of the maximum mass configuration \( C_{\text{max}} \): \( k = 2.43 - 3.38C_{\text{max}} \). In the case of SFHo and SFHo-HD, \( C_{\text{max}} = 0.3 \) and 0.23, respectively. Correspondingly, one obtains \( M_{\text{threshold}} = 2.94 \) and \( 2.61 M_{\odot} \). We can compare these values with the results that we obtain from our numerical simulations. Let us first discuss the case of SFHo-HD: from Table 2 one can notice that for \( M_{\text{tot}} > 2.48 M_{\odot} \), the remnant collapses in less than 1 ms, and \( M_{\odot} \) drops below \( m_{\odot} \). Moreover, there is almost no gravitational wave energy in the post-merger phase. Thus, from numerical simulations, we infer for SFHo-HD that \( M_{\text{threshold}} \approx 2.5 M_{\odot} \) (a few percent smaller than the estimate obtained with the empirical formula). Concerning SFHo, the largest value of \( M_{\text{tot}} \) that we have simulated is \( 2.84 M_{\odot} \), and it leads to a prompt collapse to a BH and indeed represents a good estimate of the threshold mass for the SFHo EoS (again a few percent smaller than the estimate obtained from the empirical relation). It is interesting to notice that in a more recent analysis of the threshold mass in full general relativity, a nonlinear fitting formula for the relation between \( M_{\text{threshold}} / M_{\text{TOV}} \) and \( C_{\text{max}} \) has been found; see Köppel et al. (2019). By using this new fitting formula, we obtain \( M_{\text{threshold}} = 2.86 \) and 2.52\( M_{\odot} \) for SFHo and SFHo-HD, respectively, in excellent agreement with our numerical simulations.

Finally, one can also use this fitting formula to estimate what the value of \( M_{\text{threshold}} \) would be for different hadronic EoSs or parameter sets. In particular, one can study how much the density of the formation of hyperons, which in turn determines \( M_{\text{max}}^{H} \), affects \( M_{\text{threshold}} \). Interestingly, if we artificially switch off hyperons in our calculation and keep only nucleons and delta resonances, it turns out that \( M_{\text{threshold}} \approx 2.5 M_{\odot} \) is very close to the value obtained in the presence of hyperons. In this respect, our computation of the threshold mass is mildly dependent on the (uncertain) interaction strengths of hyperons in dense matter, and it is mainly determined by the softening due to delta resonances (see Section 8 for the phenomenological implications of this result and the Appendix for an extended discussion of the hadronic EoS).

5.6. Trigger of the Phase Conversion

For \( M_{\text{tot}} < M_{\text{threshold}}^{H} \), the two-families scenario predicts that the phase conversion of hadronic matter to quark matter necessarily occurs. Simulating the process of conversion itself after the merger is numerically very challenging (one needs to couple the code used, for instance, in Herzog & Ropke 2011 and Pagliara et al. 2013 to Cactus). We can, however, estimate when the conversion should start. As explained previously, the conversion is triggered only when a significant amount of strangeness is produced during the evolution of the remnant, i.e., \( Y_{S} \gtrsim 0.2 \). In Figure 8 (left panel), we display the stripe in the mass density–temperature diagram for which the value of \( Y_{S} \) is large enough for the conversion to start. In the right panel, we show the temporal evolution of the maximum baryon density for all of the runs of the HS–HS merger. Let us first discuss the case of the run SFHo-HD 118vs118. One can see that after the first two oscillations of the remnant, the condition for quark matter nucleation is fulfilled; see also Figures 14 and 15, which allow one to visualize the regions inside the remnant where the density and temperature reach the condition for the beginning of the conversion process. The subsequent temporal evolution of the remnant would be a dramatic change of the structure of the star: within a few ms, a big part of the star is converted into quark matter, the radius of the star increases by a few km, and the heat released by the conversion heats up the star. Also, a significant amount of differential rotation can develop due to the very fast change of the total moment of inertia of the star (Pili et al. 2016). Finally, the stiffening of the EoS would stabilize the remnant with respect to the collapse. We have a similar evolution for the SFHo-HD 12 model, with the only difference being that the conditions for the phase transition are met during the second bounce instead of the third one.
A comment concerning the results for the mass ejected in these cases is in order. The oscillations of the maximum density are associated with the shock waves, which, in turn, are responsible for most of the dynamically ejected mass. Thus, even if we are not implementing the conversion to quark matter (which could modify the temporal behavior of the maximum density) in our simulations, we can assume that the results of the mass ejected by HS–HS mergers before the trigger of the phase transition (before the third peak in the SFHo-HD 118vs118 model) are fairly reliable.

Let us now discuss the cases with larger masses. For the model SFHo-HD 124vs124, which leads to a prompt collapse, the nucleation condition is met only when the collapse has already started and the formation of the quark phase cannot halt it. The intermediate case of SFHo-HD 122vs122 instead shows a first oscillation during which the maximum density reaches the threshold for nucleation, but the second oscillation already leads to the collapse of the star. In this case, it is possible that a significant part of the star is converted to quark matter, but it is quite likely that the consequent stiffening of the EoS is not strong enough to prevent the collapse.

6. Rates of the Mergers from Population Synthesis Analysis

The crucial information that we now want to provide is an estimate of the rate of events that we expect for the different types of mergers. For this purpose, we studied the evolution toward mergers of double compact objects in the two-families scenario. Specifically, we used the Startrack population synthesis code (Belczynski et al. 2002, 2008) with further updates described in Wiktorowicz et al. (2017, and references therein). As GW170817 is our main point of interest, we concentrated on systems having parameters similar to the observed ones and adapted the simulation parameters to reproduce an environment similar to that of the host galaxy (NGC 4993).

Recently, Belczynski et al. (2018) performed a study of merger rates of two NSs in the NGC 4993 environment. Here we also include QSs as possible components of the progenitor binary. Specifically, following Belczynski et al. (2018), we assumed the metallicity of the environment in which GW170817 was formed to be $Z = 0.01$, which corresponds to about 50% of the solar metallicity. The results were scaled to the observable volume of a LIGO detector for events similar to GW170817, which may be approximately represented by a sphere with a radius of $D \approx 100$ Mpc. On the basis of an Illustris simulation (Vogelsberger et al. 2014), this volume contains about $1.1 \times 10^{13} M_\odot$ of stellar mass in elliptical galaxies. The star formation history was chosen to be uniform between 3 and 7 Gyr ago (Troja et al. 2017), which is different from Belczynski et al. (2018), who used a burst-like star formation history occurring 1, 5, or 10 Gyr ago.

As far as the binary evolution is concerned, all mass transfer was assumed to be conservative; i.e., all mass lost by the donor is transferred to the accretor. Such an approach increases the rates for double compact object mergers (Chruslinska et al. 2018). On the other hand, we kept the Maxwellian distribution of natal kicks to $\sigma = 265$ km s$^{-1}$, which conforms to the observational estimates from Galactic pulsar proper motions (Hobbs et al. 2005), although lower natal kicks would increase the rates (Chruslinska et al. 2018).

Similar to Wiktorowicz et al. (2017), we adopted the two-families scenario by assuming that an HS converts into a QS when its gravitational mass reaches a maximal value, $M_{\text{max,ns}} = M_{\text{max}}^H$. During the transition, the gravitational mass of the compact object changes instantly (compared to the typical timescales of population synthesis analysis), and we update the orbital parameters accordingly.

Calculated merger rates are presented in Table 3. Two models adopt the two-families scenario with deconfinement taking place at HS masses of 1.5 and 1.6$M_\odot$ (models $M_{\text{max}}^H = 1.5$ and 1.6$M_\odot$, respectively). In the third model (one-family), which we provide for reference, all compact
For the two-families scenario, we have adopted two possible values of $M_{\text{max}}^H$ and we report the results for all of the merger combinations: HS–HS, HS–QS, and QS–QS mergers. Results for the one-family scenario are also reported for comparison. Notice that for values of $q \lesssim 0.85$, an HS–QS merger is more probable than an HS–HS merger.

Table 3

| Model  | All Mergers | GW170817-like |
|--------|-------------|---------------|
|        | HS–HS       | HS–QS | QS–QS |
|        | HS–HS       | HS–QS | QS–QS |
| $M_{\text{max}}^H = 1.5 M_\odot$ | 9.1 | 3.1 | 0.2 |
|        | 6.4 | 0.4 | 0.01 |
| $M_{\text{max}}^H = 1.6 M_\odot$ | 9.2 | 3.2 | 0.02 |
|        | 6.5 | 0.3 | 0.01 |
| One-family | 12.8 | ... | ... |
|        | 6.6 | ... | 0.3 |

Note. For the two-families scenario, we have adopted two possible values of $M_{\text{max}}^H$ and we report the results for all of the merger combinations: HS–HS, HS–QS, and QS–QS mergers. Results for the one-family scenario are also reported for comparison. Notice that for values of $q \lesssim 0.85$, an HS–QS merger is more probable than an HS–HS merger.

objects are HSs, and the deconfinement process never occurs. We note that in all models, it was assumed that all compact objects with masses above $M > 2.5M_\odot$ formed prior to the merger, collapse to a BH, and are not included in our study.

In our analysis, we compare merger rates for all systems and those corresponding to events similar to the GW170817 event. The latter were chosen as events with chirp masses in the range $M_{\text{chirp}} = (1.188 \pm 0.1)M_\odot$ (the $\pm 0.1M_\odot$ was chosen arbitrarily, but the specific adopted value of the range has a negligible effect on the conclusions) and a mass ratio of $q > 0.7$ ($q = m_2 / m_1 > 0.7$, where $m_2 \leq m_1$).

First, when considering all mergers, it is noticeable that the rate of HS–QS is not strongly suppressed in respect to the rate of HS–HS; they constitute roughly one-fourth of the total rate.

Second, concerning the GW170817-like events, different types of mergers are allowed depending on the value of the mass ratio $q$. In Table 3 we compare the merger rates for the entire range of mass ratios ($q > 0.7$) and, separately, for lower mass ratios ($0.7 < q < 0.85$). The latter was motivated by an analysis of the GW170817 event (Abbott et al. 2018) indicating that a mass ratio $q \approx 0.85$ is highly compatible with the data. Results for both ranges of mass ratios are significantly different. The entire range ($q > 0.7$) is dominated by high mass ratio systems that are mostly HS–HS, because for $q \approx 1$ and $M_{\text{chirp}} \approx 1.188$, the components’ masses can barely enter the QS mass range $M_{\text{min}}^Q > 1.37$ or $1.46M_\odot$ (for $M_{\text{max}}^H = 1.5$ and $1.6 M_\odot$, respectively). On the other hand, for low mass ratios, one of the stars usually becomes more massive than $M_{\text{min}}^Q$, thus, according to the two-families scenarios, it converts into a QS. As a result, although the mergers within the entire mass ratio range ($q > 0.7$) show a preference for HS–HS mergers, the ones with a lower mass ratio ($0.7 < q < 0.85$) are typically HS–QS mergers. However, the former give much higher merger rates ($\sim 6.8$ kyr$^{-1}$) than the latter ($\sim 0.2$–0.3 kyr$^{-1}$).

It is also important to note that the rates of QS–QS mergers are in general very much suppressed, reaching at most a few percent of the total expected events involving at least one QS in the merger, a result in agreement with the findings of Wiktorowicz et al. (2017). In particular, the probability that GW170817 was due to a QS–QS merger is totally negligible.

Finally, we note that all of the merger rates obtained from population studies are significantly below the one estimated from the GW170817 event (by more than 2 orders of magnitude), as previously shown by Belczynski et al. (2018). It may be a result of, e.g., low observational statistics (just one event), poorly understood binary evolution phases (e.g.,

common envelope survival of low mass ratio binaries), or the existence of an unknown evolutionary scenario.

In the following, we discuss the interpretations of GW170817 in terms of the different types of mergers and their probability.

### 6.1. HS–HS Binary

Our results show that, provided the entire possible range of mass ratio values ($q > 0.7$) is taken into account, in the two-families scenario, GW170817 is most likely a merger of two HSs. Indeed, while the observationally estimated chirp mass ($M_{\text{chirp}} \approx 1.188M_\odot$) and mass ratio ($q > 0.7$) limit the mass of the primary to $(1.37-1.64)M_\odot$ and the secondary to $(1.17-1.37)M_\odot$, a result of population synthesis analysis is that mass ratios $q \sim 1$ are favored, and in equal-mass binaries having $M \sim 1.37M_\odot$, QSs are very unlikely to occur.

In a typical evolution, which produces a double HS merger in the two-families scenario, the masses of the binary components on the zero-age main sequence (ZAMS) are 9.7 and 8.6$M_\odot$, and the separation is moderate ($a \approx 340R_\odot$). The primary, i.e., the heavier star on ZAMS, evolves faster and, after 27 Myr, expands as a Hertzsprung gap star, fills its Roche lobe, and commences mass transfer onto its companion. The mass transfer lasts $\sim 20,000$ yr. Afterward, the primary is ripped off its hydrogen envelope and left as a $2.1M_\odot$ helium star. The companion accretes all of the transferred mass (as assumed in our computations) and reaches a mass of 16$M_\odot$.

After an additional 4 Myr, the primary explodes as a supernova (SN) and forms a $1.26M_\odot$ compact object. At that moment, the separation is still high ($a \approx 1900R_\odot$). The secondary evolves and, after an additional 4 Myr, expands and fills its Roche lobe as an AGB star. This time, the mass ratio is too extreme ($q \approx 0.24$), and the mass transfer is unstable, leading to a common envelope. As a consequence, the separation shrinks to $a \approx 29R_\odot$, and the secondary becomes a helium star. Shortly after, it forms a compact object in an SN explosion. The natal kick slightly enlarges the orbit to $a \approx 95R_\odot$, but the eccentricity, $e \approx 0.98$, is large enough to allow for a merger within $\sim 5$ Gyr.

### 6.2. HS–QS Binary

In the case of low mass ratios ($0.7 < q < 0.85$), the most probable scenario involves the merger of an HS and a QS, because the most massive component will typically be massive enough to become a QS. This result is true in general, and it also holds for GW170817-like events. Notice, though, that for
GW170817-like events, the total merger rate for low mass ratios (≈0.2–0.3 kyr⁻¹) is more than an order of magnitude smaller than for the wider range (i.e., q > 0.7; ≈6.8 kyr⁻¹).

In this situation, the progenitor binary on the ZAMS typically consists of 15.9 and 15.6M☉ stars on an eccentric orbit (e ≈ 0.8) with a separation of approximately a ≈ 410R☉. The primary fills its Roche lobe first and commences mass transfer onto the counterpart. The outcome is a mass reversal. The primary becomes an ∼4.4M☉ helium star, whereas the secondary is now a 26.5M☉ main-sequence star. About 1 Myr later, the primary forms an ∼1.2M☉ NS in an SN explosion. The secondary quickly evolves, expands, and fills its Roche lobe as a Hertzsprung gap star. The donor is much heavier than the accretor, so the mass flow becomes unstable and results in a common envelope. The binary survives the phase, but separation is significantly smaller (a ≈ 2R☉). The secondary loses most of its mass and becomes an ∼8M☉ helium star that, after a short evolution (∼1 Myr), forms a compact object in an SN and becomes an ∼1.5M☉ QS because it was too heavy to form an HS. The separation grows to ∼4.7R☉ due to natal kick, but significant eccentricity allows for a merger within the next ∼4.6 Gyr.

6.3. QS–QS Binary

As already discussed above, a GW170817-like merger is very unlikely to be of the QS–QS type. More precisely, GW170817 cannot be a QS–QS, as is immediately evident from Figure 3. The GW170817-like mergers can be marginally QS–QS for q ∼ 1, where they fit in the “All-comb.” region of Figure 3.

7. GW170817 and Its Interpretation as an HS–QS Merger

In the case of GW170817, the total mass is \( M_{\text{tot}}^{170817} = 2.74 \pm 0.04 M_\odot \), and, as explained before, within the two-families scenario, this event can be interpreted only as due to the merger of an HS with a QS. Indeed, \( M_{\text{threshold}} < M_{\text{tot}}^{170817} \), and therefore an HS–HS binary would promptly collapse to a BH. However, a direct collapse is excluded by the observation of a sGRB \( \sim 2 \) s after the merger. Also, as discussed previously, it cannot be interpreted as a merger of two QSs.

A first question concerns the probability of detecting such an HS–QS merger event. From the binary evolution point of view, the population synthesis analysis has shown that the delay time, which is ∼4.6 Gyr, is consistent with the lack of recent star formation in the host galaxy.

Concerning the phenomenology of GW170817, without detailed numerical simulations of the merger of an HS with a QS, it is difficult to firmly establish whether we can explain all the features of the observed signals in our scenario. We can, however, verify whether our model satisfies the upper and lower limits on \( \Lambda \). The lower limit on \( \Lambda \) in particular, empirically found in Radice et al. (2018b), stems from the request for a description of the kilonova signal AT 2017gfo (this constraint was recently criticized in Kiuchi et al. 2019).

In Figure 16, we show the tidal deformabilities of the two components of the binary and \( \Lambda \) for the chirp mass of GW170817 for the one-family and two-families scenarios (for the cases of HS–HS, HS–QS, and QS–QS mergers14). Both the upper limit on \( \Lambda \), as obtained by the gravitational wave’s signal, and the lower limit, as inferred from the kilonova analysis (Radice et al. 2018b), can be fulfilled within the two-families scenario in the case of an HS–QS system (see also Burgio et al. 2018). Notice that the small value of \( \Lambda \) for the HS component is due mainly to the softening associated with the appearance of delta resonances.

We observe that the case of HS–HS is not compatible with the constraints on \( \Lambda \) because it provides a too-small value for it (Radice et al. 2018b; Kiuchi et al. 2019). Thus, also from the constraints on \( \Lambda \), we can rule out the possibility that GW170817 was due to an HS–HS merger. We need to discuss now how the electromagnetic counterparts of GW170817 would be produced in our scheme.

Concerning the sGRB, we can assume that it has been produced by the disk accreting onto the BH in agreement with the classification discussed in Section 4. Indeed, for GRB 170817A, there is no clear indication of the existence of an extended emission after the prompt signal. Concerning the kilonova, a crucial issue is connected with the fate of the quark matter ejected from the QS during the merger. There are a few studies indicating that quark matter evaporates into nucleons and that this process can be so efficient that only strangelets with baryon numbers much larger than \( 10^{40} \) can survive (Alcock & Farhi 1985; Madsen et al. 1986; N. Bucciantini et al. 2019, in preparation). Most of the matter ejected from the QS will therefore rapidly evaporate into nucleons and can contribute to the kilonova signal. If this idea is correct, the kilonova produced after an HS–QS merger would not show significant differences, concerning the type of matter ejected, with respect to a kilonova produced after an HS–HS merger. Actually, the process of matter ejection can be very efficient, and it opens up the possibility of explaining AT 2017gfo as an HS–QS merger: in the case of AT 2017gfo, the radii of the two compact objects are both rather small, the system is asymmetric, and the threshold mass is large, as discussed in Section 3. These are exactly the requests posed in Kiuchi et al. (2019), and, as discussed in that paper, there are no examples.

14 The QS–QS case is shown for comparison, but it cannot be realized because it violates the limit on \( M_{\text{min}}^{\text{G}} \).

Figure 16. Tidal deformabilities for the components of a binary system (one-family and two-families scenarios) for a value of the chirp mass equal to the one of GW170817. For the one-family scenario, we use the SFHo EoS. For the two-families scenario, we use the SFHo-HD and QS EoSs, and we consider the possibility of HS–HS, HS–QS, or QS–QS systems.
of EoSs based on a single-family scenario able to satisfy all of those requests. Finally, we note that while the scenario of a delayed collapse (a few tens of ms after the merger) is adopted by most of the literature on GW170817, there are some studies indicating that the merger remnant could be a supramassive or even a stable star (Ai et al. 2018; Li et al. 2018b) and that it could have been active for a much longer timescale. In van Putten & Della Valle (2018), a post-merger gravitational wave signal has been reported with an associated gravitational wave energy lower than the sensitivity estimates of the LIGO-VIRGO collaboration. This signal would have lasted a few seconds, with an initial frequency of about 700 Hz. Another possible signal of a long-lived remnant comes from the analysis of Piro et al. (2019), in which was found a low-significance temporal feature in the X-ray spectrum ~150 days after the merger that is consistent with a sudden reactivation of the central compact star. If these analyses are correct, they can have very important implications for the EoS of dense matter that we will address in a forthcoming study.

8. Discussion and Conclusions

We have explored the phenomenological consequences of the existence of two families of compact stars on the observations of mergers. A specific feature of the two-families scenario is the possibility that at fixed values of the chirp mass and mass asymmetry, the merging binary can be composed by two HSs, two QSs, or an HS and a QS. Each of these possibilities has its own specific signatures.

Let us first consider systems similar to the one associated with GW170817. Our population synthesis analysis basically excludes the possibility that such binary systems are made of two QSs (notice that this result depends only on the adopted value for $M_{\text{max}}$, which is rather tightly determined in our scenario). They can be formed either by two HSs or by an HS and a QS.

If we look to the whole range of values for the mass ratio $q$ of the two stars, the rate of HS–HS merger events is a factor of (16–20) larger than the rate of HS–QS mergers, while for a small mass ratio ($0.7 < q \lesssim 0.85$), HS–QS systems are about two to six times more probable than HS–HS systems. While in the HS–HS case, we would predict a prompt collapse (no sGRB and a very faint kilonova), in the HS–QS case, a hypermassive remnant is more likely to form (with associated electromagnetic signals similar to GRB 170817A and AT 2017gfo). Notice that within the one-family scenario, the total mass associated with GW170817 can be regarded as a lower limit for the threshold mass, since there are clear indications that a direct collapse to a BH did not take place. In the two-families scenario, the threshold mass instead depends on the type of merger, and it is dramatically different for an HS–HS and an HS–QS merger. In particular, for the HS–HS case, the threshold mass, as found in Section 5.5, is of the order of $2.5M_\odot$ and thus significantly smaller than the mass associated with GW170817.

The smoking gun of our scenario would be just a single detection of a source of gravitational waves with a total mass smaller than the one of GW170817 but lacking a significant electromagnetic counterpart (indeed, it would be interpreted as an HS–HS merger). This possible signature is valid for values of the total mass in the range (2.48–2.74)$M_\odot$, i.e., between the threshold mass for the merger of two HSs and the mass of the system associated with GW170817.

Another promising way to test the two-families scenario is related to the observation of binaries with low values of $M_{\text{tot}}$. Let us concentrate on $M_{\text{tot}} = 2.4M_\odot$ with two stars of $1.2M_\odot$. A first difference between the two-families and the one-family scenarios concerns the value of $\Lambda$, which can be significantly smaller (almost a factor of 2 for our specific choice of the hadronic EoS) for a merger of HS–HS with respect to the merger of two NSs (see the values for SFHo-HD and SFHo in Table 1). This is due to the appearance of delta resonances in hadronic matter at densities of about twice the saturation density. Accurate future measurements of the gravitational wave’s signal associated with the inspiral phase should put interesting upper limits on $\Lambda$ and therefore test the possibility of the merger of very compact stellar objects, i.e., HSs in our scheme.

Additionally, if observations of the post-merger signal will be feasible in future experiments (most likely the third-generation detectors), one can test the two-families scenario through the measurement of the frequency of the $f_2$ mode, which, for an HS–HS merger, is of the order of 1 kHz larger than the case of an NS–NS merger during the first milliseconds of the life of the remnant. The subsequent formation of quark matter, which in our scenario corresponds to a stiffening of the EoS, should then decrease the value of $f_2$ (see Bauswein et al. 2016 for a preliminary analysis of this signature). It is interesting to note that in recent studies discussing a phase transition to quark matter in the post-merger, the appearance of quark matter reduces the lifetime of the remnant and leads to a shift of $f_2$ to frequencies larger than those of the “progenitor” made only of nucleonic matter (Bauswein et al. 2019; Most et al. 2019).

Let us discuss the possible signatures associated with the kilonova. Its signal depends strongly on the amount of mass dynamically ejected by the merger and in the disk. One can notice from Table 2 that for an HS–HS merger, both of these masses are of the order of $0.01M_\odot$. Instead, in the case of an NS–NS merger (in the one-family scenario), the dynamically ejected mass is significantly smaller than the mass of the disk. Qualitatively, we would then expect a significant difference between the resulting kilonova signals: within the two-families scenario (in the case of an HS–HS merger), the intermediate-opacity purple component is strongly suppressed with respect to the components associated with the dynamically ejected matter (Perego et al. 2017).

Finally, it is interesting to notice that the suggestion of the two-families scenario originated from the requirement of explaining the observations of stars having large masses and the possible existence of stars with small radii. The same problem presents itself again when trying to model the kilonova AT 2017gfo: as suggested in Siegel (2019), the presence of a strong shock-heated component in the ejecta would favor radii smaller than (11–12) km. Such small radii can be easily accommodated in the two-families scenario but are basically ruled out within the one-family scenario.

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15 http://universeathome.pl
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Appendix

In this Appendix, we will summarize the results of the papers published during the last 5 yr leading to a precise definition of the two-families scenario. The aim of this section is to justify the selection of the two EoSs we use for the hadronic and quark families and to explain why we consider those equations as representative of the entire class of EoSs that can be used in that scenario. One important point is that our approach is mainly phenomenological, based on the attempt to interpret the existing data (although controversial) and provide predictions directly related to the interpretation of those data.

The main motivation for proposing the two-families scenario in 2014 (Drago et al. 2014a) was the need to reconcile the existence of very massive stars and the possible existence of very compact objects. In particular, several studies on low-mass X-ray binaries have suggested values for the radii significantly smaller than ∼12 km (see Figure 17 where we display the observational constraints indicated by a couple of very recent analyses; for other examples, see Drago et al. 2016b). While the values of the masses of those objects are not known (but masses smaller than about 1.5M⊙ are favored; d’Etivaux et al. 2019), it is unlikely that all of those objects have masses much larger than the canonical value of 1.4M⊙. Therefore, these very recent analyses (as many others in the recent past) suggest that R1.4 could be significantly smaller than 12 km. If we assume that those estimates for the radii are correct and interpret those stars as HSs, we need to tune the parameters of our hadronic model in order to explain such small radii. Typically, small radii imply large central densities and make the opening of softening channels, such as the production of delta resonances (Schürhoff et al. 2010) and/or hyperons (Chatterjee & Vidaña 2016), very likely. It is therefore very simple and direct to relate small radii to the appearance of new degrees of freedom: since, in the two-families scenario, the hadronic branch needs not support very massive stars, one need not introduce ad hoc repulsive mechanisms that from one side would avoid the production of, e.g., hyperons but on the other side would not allow for small values of R1.4 (an example of that mechanism is discussed in, e.g., Lonardoni et al. 2015).

An important point of our scenario is that the nucleation of quark matter takes place only when a sizable amount of strangeness is present in the HS. An early appearance of hyperons would therefore imply an early transition to quark matter, implying that stars having a mass larger than about 1.4M⊙ would convert to QSs. This scenario would most likely be ruled out by the phenomenology associated with, for instance, magnetar X-bursts. The early formation of delta resonances solves this problem because it shifts the formation of hyperons to larger densities, as found in Drago et al. (2014b) and confirmed in, e.g., Li et al. (2018a) and Li & Sedrakian (2019). Another issue associated with the late appearance of hyperons concerns the cooling: as was shown in Prakash et al. (1992), if, for a certain stellar configuration, the direct nucleon URCA mechanism is forbidden, then the direct URCA process for ∆ is also forbidden. Thus, we expect stars containing only nucleons and ∆s to behave as slow coolers. On the other hand, stars close to the end of the hadronic branch (with central densities above ∼5n0) contain enough ∆s that the direct URCA is active and would instead cool rapidly. Interestingly, the phenomenology of the cooling of compact stars indicates that, at least for some cases, fast cooling is needed to explain the data (Brown et al. 2018).

Let us first discuss the delta resonances. The couplings between them and mesons (within the relativistic mean field approach that we use here) because of SU(2) symmetry must be of the same order of magnitude as the couplings between nucleons and mesons. In particular, in Drago et al. (2014b), it was realized that the properties of the EoS are determined by the relative magnitude between the couplings of deltas and scalar mesons and of deltas and vector mesons. Thus, we can fix the couplings with vector mesons to be equal to the couplings between nucleons and vector mesons and regard the coupling with scalar mesons ξΔ to be a free parameter. Actually, electron–nucleus scattering data constrain ξΔ to vary in the range 1–1.15. With the corresponding EoSs, we obtain the dashed mass–radius curves shown in Figure 17 for the extremes of the allowed range for ξΔ.

Let us now discuss the effect of the appearance of hyperons. As in the case of delta resonances, their effects on the EoS in our framework are determined by the couplings with mesons. Moreover, one should extend the discussion to SU(3) and consider the full scalar and vector mesons nonets. There are clearly too many free parameters and only a few experimental data: one can use SU(3) symmetry as a guideline (as assumed here), but violations are possible; see Weissenborn et al. (2012).
Interestingly, the formation of hyperons is crucial to fix the value of $M_{\text{max}}^H$, which cannot be very small, as discussed above, but also cannot be too large. For instance, a value of $M_{\text{max}}^H \sim 1.8M_\odot$ would lead to a minimum mass of QSn, $M_{\text{min}}^H \sim M_{\text{max}}^H - 0.1M_\odot \sim 1.7M_\odot$. Therefore, only stars more massive than $1.7M_\odot$ could have large radii.

On the other hand, there are some indications of stars with a mass close to 1.5$M_\odot$ and a large radius: of the order of 13 km for the cases with and without hyperons, respectively. Instead, by setting $x_{\delta,\Delta} = 1$ (see black lines in Figure 17), we obtain $M_{\text{threshold}}^H = 2.6$ and $2.68M_\odot$ for the cases with deltas and hyperons and only deltas, respectively. Those values for the threshold mass are clearly larger with respect to the case of SFHo-HD but still smaller than the threshold mass predicted by the SFHo EoS and the mass associated with GW170817.

In this respect, our prediction for the cases of prompt collapses would be only mildly dependent on the threshold for the appearance of hyperons, and it would be mainly determined by the effect of introducing the delta resonances in the hadronic EoS (with a specific value of the coupling).

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