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Neutron star matter at finite temperature and implications for neutron star mergers

Hristijan Kochankovski



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PhD Thesis

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Neutron star matter at finite temperature and implications for neutron star mergers

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In honor of my grandpa Jole

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Matèria d'estrelles de neutrons a temperatura finita i implicacions per a les fusions d'estrelles de neutrons

Resum

Paraules clau: Estrelles de neutrons, Hiperons, Equació d'Estat, Temperatura finita, Fusions binàries d'estrelles de neutrons, Ones gravitacionales

L'aparició d'hiperons en matèria ultra-densa i calenta com la que es troba a les estrelles de neutrons (NS) és el tema central que s'investiga en aquesta tesi. Abordem el tema des de diferents perspectives. En primer lloc, estudiem un formalisme teòric en el qual els hiperons es poden incloure a nivell de l'equació d'estat (EoS). En segon lloc, estudiem les diferents empremtes que deixen en els observables de les estrelles a partir d'estrelles aïllades i fredes i d'esdeveniments de fusió de sistemes binaris d'estrelles de neutrons (BNS).

Desenvolupem una EoS en el marc del camp mitjà relativista, anomenada FSU2H*, que incorpora explícitament els hiperons com a grau de llibertat per descriure la matèria bariònica homogènia en condicions de densitat, fracció de càrrega i temperatura arbitràries. El model es basa en l'anterior esquema FSU2H que es va desenvolupar principalment per descriure la matèria freda d'estrelles de neutrons, tenint en compte les propietats de saturació de la matèria nuclear, les propietats dels nuclis finits, així com les restriccions astrofísiques dels púlsars massius. Els paràmetres del model FSU2H* s'ajusten per tenir en compte els resultats més recents sobre els potencials dels hiperons en matèria nuclear simètrica i l'enllaç $\Lambda\Lambda$ en matèria Λ . El model també s'estén a temperatures finites, cosa que permet el càlcul de l'EoS per a matèria densa i calenta amb diverses composicions, cobrint les condicions rellevants per a les explosions de supernoves, les proto-estrelles de neutrons (PNS) i les fusions de BNS.

Una part d'aquesta tesi també es dedica a explorar sistemàticament com les incerteses en les dades hipernuclears existents influeixen en l'EoS hiperònica a temperatura finita. Desenvolupem dues parametritzacions addicionals, FSU2H*L i FSU2H*U, que es basen en el nostre model nominal FSU2H*. Aquestes parametritzacions cobreixen de manera efectiva el rang d'incertesa dels potencials dels hiperons, representant els límits inferiors

(L) i superiors (U), cosa que condueix a EoS hiperòniques més toves i més rígides, encara consistents amb les dades experimentals extretes dels hipernuclis disponibles. Mostrem explícitament que la variació del model es propaga a una incertesa de les nostres prediccions tant en la composició de la matèria com en la mateixa EoS. La predicció de la pressió en matèria ultra-densa es veu particularment afectada, mentre que l'entropia per partícula i l'energia per partícula mostren menys sensibilitat a les variacions dels potencials dels hiperons dins de la matèria.

Es presta una atenció especial a la matèria de NS en equilibri beta sense neutrins, ja que aquest escenari és important per a la comprensió de les estrelles aïllades evolucionades i imita de prop les fases finals de l'evolució de les PNS. Inspeccionem el comportament tèrmic de l'EoS amb el nostre model nominal, a través de l'índex tèrmic. Les nostres troballes mostren que, si bé les variacions de l'índex tèrmic amb la temperatura i la densitat són suaus per a la matèria purament nucleònica, aquest no és el cas per a la matèria hipernuclear. L'aparició d'hiperons indueix una caiguda de l'índex tèrmic que és més forta a temperatures més baixes $T < 20$ MeV i fins i tot pot conduir a valors negatius. A temperatures més altes, aquesta caiguda esdevé més suau, però s'estén sobre un rang més ampli en l'espai de densitat. Aquests resultats poden tenir implicacions significatives per a les simulacions relativistes de supernoves i fusions d'estrelles de neutrons, que convencionalment es basen en una EoS freda complementada amb un comportament tèrmic que assumeix un valor constant de l'índex tèrmic (sovint anomenat la llei Γ). El nostre treball demostra que aquesta aproximació de valor constant és inadequada per reproduir amb precisió els veritables efectes tèrmics, especialment quan els hiperons esdevenen abundants en les condicions de matèria calenta i densa que es troben en aquests esdeveniments astrofísics. Per fer que els nostres resultats siguin pràcticament útils, també hem proporcionat una parametrització simple de l'índex tèrmic que es pot utilitzar en simulacions relativistes.

També investiguem com la incertesa en les interaccions dels hiperons es tradueix en les propietats globals predites de les estrelles, tant a temperatura zero com a temperatura finita. Es realitzen càlculs per a la relació massa-radi de les estrelles, la seva deformabilitat de marea i el moment d'inèrcia. Per als càlculs a temperatura zero en equilibri beta, comparem els resultats obtinguts amb les restriccions que tenim de les mesures astrofísiques. Mostrem que les incerteses hiperòniques van influeixen en la massa màxima de NS freda que l'EoS pot predir, amb $0.14M_{\odot}$ de diferència entre la predicció de la nostra parametrització més rígida FSU2H*U i la parametrització més tova FSU2H*L. A més, per a la majoria d'estrelles massives, la diferència del radi predit per a una estrella amb massa fixa pot arribar fins al 10%. No obstant això, l'impacte de les incerteses en els radis corresponents de les estrelles de massa màxima és negligible, ja que no supera els 0.2 km. Vam obtenir resultats similars per a la deformabilitat de marea i també per al moment d'inèrcia. La deformabilitat de marea a altes masses de NS es veu especialment

afectada per les incerteses, ja que les prediccions poden ser diferents per un factor de ≈ 2.8 .

Els càlculs a temperatura finita es realitzen utilitzant perfils de temperatura i leptons per a les estrelles motivats per les condicions presents en l'evolució de les PNS tant en la seva evolució primerenca (al voltant d'1 s després de l'explosió de l'estrella) com en la seva fase tardana d'evolució (desenes de segons després de l'explosió de l'estrella). Trobem que les estrelles durant la seva evolució primerenca tenen una compacitat inferior, i connectem aquest fet amb la rigidesa de la EoS de baixa densitat en matèria calenta. La influència de la incertesa hiperònica en els observables 'es anàloga a la de temperatura zero. No obstant això, a causa de la major fracció de leptons en la fase primerenca, l'aparició d'hiperons es veu obstaculitzada en aquesta evolució, cosa que condueix a un nombre quàntic d'estranyesa acumulat inferior per a aquests casos, en comparació amb els casos en què es considera que la matèria està en equilibri β^- sense neutrins. En conseqüència, l'impacte de les incerteses és més gran en els observables calculats per a les estrelles en les seves fases d'evolució tardana.

La segona part de la tesi es dedica a una anàlisi que va més enllà d'una única descripció del model de la matèria. Utilitzant 14 EoS hiperòniques diferents, inspeccionem les seves propietats a temperatura finita. Confirmem que la caiguda de la pressió tèrmica quan els hiperons apareixen en la matèria és present independentment del model o del marc utilitzat. Ho connectem amb l'aparició d'hiperons tèrmics en la matèria. A més, mostrem que la matèria hipernuclear en general mostra una capacitat calorífica més alta que la matèria nuclear ordinària.

Els resultats van motivar una extensa campanya de simulació que es va realitzar amb les esmentades EoS hiperòniques i també amb 23 EoS purament nucleòniques. L'objectiu de la campanya era identificar correctament possibles empremtes hiperòniques en els observables de la fusió de BNS i discutir el seu efecte en l'evolució de diferents quantitats. Al mateix temps, la campanya de simulació ens va permetre estudiar com els diferents comportaments tèrmics afecten el resultat de les simulacions.

Dels resultats obtinguts, observem un augment petit (al voltant de pocs per cent) però sistemàtic de la freqüència dominant d'ona gravitatòria (GW) post-fusió quan els hiperons tèrmics estan presents en la matèria romanent respecte als casos en què no hi són presents. El desplaçament de la freqüència anti-correlaciona amb l'índex tèrmic de la matèria, cosa que confirma el fet que el reblaniment de l'EoS de temperatura finita induït pels hiperons tèrmics desencadena aquest fenomen. Corroborem que l'augment característic no depèn de les condicions inicials de les nostres simulacions; l'efecte és present en les nostres simulacions nominals (binàries simètriques amb una massa total de $2.8M_{\odot}$), en les nostres simulacions binàries asimètriques mantenint la mateixa massa total de $2.8M_{\odot}$.

però amb una asimetria entre les masses de les estrelles de $M_1/M_2 = 0.8$ i en les nostres simulacions simètriques més massives.

També explorem la influència dels hiperons en altres característiques dels espectres de GW. Mostrem explícitament que el pic secundari f_{2-0} pateix un augment similar quan els hiperons tèrmics estan presents en la matèria, mentre que el pic secundari f_{spiral} no es veu afectat per l'aparició d'hiperons. Relacionem el diferent comportament qualitatiu de les dues característiques amb els diferents mecanismes en què es produeixen.

També estudiem la temperatura mitjana en els romanents. La capacitat calorífica específica més alta de la matèria hipernuclear induïx temperatures més baixes en els romanents en comparació amb els casos en què només hi havia matèria nuclear.

També es mostren els resultats per a la massa ejectada dinàmica i la massa llinar per al col·lapse a forat negre immediat. En considerar models hiperònics, la massa ejectada dinàmica durant les fusions sembla ser més gran que en les EoS nucleòniques, per a casos en què ambdós produeixen propietats de NS fredes similars. A més, les EoS hiperòniques condueixen a una massa llinar inferior per a la formació immediata de forat negre (aproximadament $0.05M_{\odot}$) en comparació amb les seves contraparts nucleòniques, assumint els mateixos paràmetres estel·lars de NS freda.

Abstract

This PhD thesis investigates the properties of neutron star (NS) matter, with a primary focus on the appearance and effects of hyperons at finite temperature. The work aims to contribute to resolving the "hyperon puzzle" in a two-fold way. On the one hand, in the thesis a theoretical framework necessary to describe hot and dense baryon matter is developed. The relativistic mean field approach is used, as it naturally incorporates relativistic effects, explicitly includes meson degrees of freedom, and can be generalized to finite temperature. The FSU2H* equation of state (EoS) is developed, which integrates the experimental data on hyperon potentials in symmetric nuclear matter and $\Lambda\Lambda$ bond energy to ensure consistency with nuclear and hypernuclear physics constraints.

Furthermore, two extreme parametrizations of the models are constructed so as to cover the uncertainties in the hyperon-meson couplings arising from the scarce experimental data. The two additional parametrizations, FSU2H*L and FSU2H*U, represent the EoSs that satisfy the lower and upper limits of these uncertainties, respectively. It is demonstrated that the models are encompassing the uncertainty range in the total hyperonic content and in the EoS itself in the conditions relevant for supernovae, proto-neutron stars (PNSs), and binary neutron star (BNS) mergers.

It is shown that the pressure is especially sensitive to the variation of the hyperon-meson couplings, which in the densest part of the stars shows a 10% difference between the different parametrizations. Other thermodynamical quantities like the energy per particle and entropy per particle only show small variations.

With the constructed parametrizations, several stellar observables are computed both in cold and hot stars. For cold NS properties, hyperonic uncertainties significantly impact the maximum NS mass (a difference of $0.14 M_{\odot}$ between FSU2H*U and FSU2H*L) and the radius of massive stars (up to 10% difference for a fixed mass). The tidal deformability of very massive cold neutron stars is particularly affected, with predictions differing by a factor of approximately 2.8 at high NS masses due to these uncertainties. For hot NS properties, calculations for stars resembling phases of the PNS evolution demonstrate that higher temperatures at lower densities lead to lower compactness. The influence of hyperonic uncertainties on hot star observables mirrors that for cold stars, with a greater

impact observed in later PNS evolution phases due to lower lepton fractions promoting hyperon appearance.

On the other hand, the thesis also searches for model-independent signatures of exotic baryonic matter. Inspecting 14 different hyperonic EoSs and 23 nucleonic EoSs, it is demonstrated that the appearance of hyperons induces a characteristic drop in the thermal pressure and, consequently, the thermal index. This behavior is attributed to the redistribution of degeneracy pressure as new, less degenerate heavy baryons are thermally populated. Furthermore, hypernuclear matter generally exhibits a higher specific heat at constant pressure compared to purely nucleonic matter, indicating an enhanced capacity to store thermal energy due to the opening of new degrees of freedom.

With an extensive simulation campaign using all EoSs at our disposal, the imprints of hyperons on BNS merger observables and remnant evolution quantities are studied. The main finding is a systematic positive shift (2-4%) in the dominant post-merger GW frequency, which is identified when hyperonic models are compared to reference simulations that assume an ideal nucleonic thermal behavior. This shift is anti-correlated with the average thermal index in matter and is robust across variations in total binary mass and mass asymmetry. The shift, estimated to be between 50 and 150 Hz, is potentially detectable by future gravitational-wave observatories such as the Einstein Telescope or Cosmic Explorer.

In addition, the secondary GW peak f_{2-0} , exhibits similar positive shifts, while f_{spiral} remains largely unaffected by the thermal behavior. The numerical robustness of the identified GW shifts is confirmed through comparisons with different numerical schemes and resolution studies.

Some non-spectral observables are studied too. Hyperonic models predict larger dynamical mass ejecta (up to a factor of two) compared to purely nucleonic models with similar cold NS properties. The presence of hyperons also leads to a reduction in the threshold mass for prompt black hole collapse by approximately $0.05 M_{\odot}$, attributed to the thermal softening of the EoS.

Glossary

BHF Brueckner-Hartree-Fock. 17

BNS Binary Neutron Star. 5

CCSN Core Collapse Supernova. 34

CE Cosmic Explorer. 14

CFA Conformal Flatness Approximation. 71

CMF Chiral Mean Field. 62

DFT Density Functional Theory. 60

EoS Equation Of State. 4

ET Einstein Telescope. 14

GR General Relativity. 12

GWs Gravitational Waves. 13

J-PARC The Japan Proton Accelerator Research Complex. 27

KAGRA Kamioka Gravitational Wave Detector. 13

KEK The High Energy Accelerator Research Organization. 27

LIGO The Laser Interferometer Gravitational-Wave Observatory. 13

NICER Neutron Star Interior Composition Explorer. 12

NSs Neutron Stars. 4

PNM Pure Neutron Matter. 26

PNS Proto-Neutron Star. 7

QCD Quantum Chromodynamics. 17

QMC Quark-Meson-Coupling. 62

RMF Relativistic Mean Field. 5

sGRBs Short Gamma-Ray Bursts. 14

SKA Square Kilometer Array. 13

SNM Symmetric Nuclear Matter. 26

SPH Smoothed Particle Hydrodynamics. 71

TOV Tolman-Oppenheimer-Volkoff. 44

YN Hyperon-Nucleon. 18

YY Hyperon-Hyperon. 18

List of publications

The following articles have been published as a result of the work presented in the thesis:

- **H. Kochankovski**, À. Ramos, L. Tolos, "**Equation of state for hot hyperonic neutron star matter**", *Monthly Notices of the Royal Astronomical Society*, **517** (2022), 1, 507-517; **518** (2023), 4, 6376.
- **H. Kochankovski**, À. Ramos, L. Tolos, "**Hyperonic uncertainties in neutron stars, mergers, and supernovae**", *Monthly Notices of the Royal Astronomical Society*, **528** (2024), 2, 2629–2642.
- S. Blacker, **H. Kochankovski**, A. Bauswein, À. Ramos, L. Tolos, "**Thermal behavior as indicator for hyperons in binary neutron star merger remnants**", *Physics Review D*, **109** (2024), 043015.
- **H. Kochankovski**, G. Lioutas, S. Blacker, A. Bauswein, À. Ramos, L. Tolos, "**Impact of hyperons on neutron star mergers: Gravitational waves, mass ejection and black hole formation**", *Physics Review D*, **112** (2025), 023014.

During the time as a PhD student at the University of Barcelona, the following articles were also published:

- **H. Kochankovski**, À. Ramos, I. Vidaña, "**An analytic parametrization of the hypernuclear matter equation of state**", *European Physics Journal A*, **58**, 31 (2022).
- S. Gjorgievska, **H. Kochankovski**, K. Stankovic, L. Barandovski, "**Revision of the semi-empirical mass formula coefficients by using the ame2020 database**", *Nuclear Engineering and Design*, **426** (2024), 113403.
- T. Fischer, J. M. Camalich, **H. Kochankovski**, L. Tolos, "**Hyperons during proto-neutron star deleptonization and the emission of dark flavoured particles**", *Journal of Cosmology and Astroparticle Physics*, **01** (2025), 061.

- A. Bauswein, A. Nikolaidis, G. Lioutas, **H. Kochankovski**, P. Char, C. Mondal, M. Oertel, L. Tolos, N. Chamel, S. Goriely, "**Stellar properties indicating the presence of hyperons in neutron stars**", *Submitted*.
- **H. Kochankovski**, À. Ramos, L. Tolos, "**Momentum-fragmented hyperons in cold neutron stars**", *Submitted*.

The following conference proceedings were also published:

- **H. Kochankovski**, À. Ramos, L. Tolos, "**Finite-temperature equation of state with hyperons**", *EPJ Web of Conferences*, **274** (2022), 07004.
- **H. Kochankovski**, À. Ramos, L. Tolos, "**Hyperonic equation of state at finite temperatures for neutron stars**", *EPJ Web of Conferences*, **271** (2022), 09005.
- **H. Kochankovski**, À. Ramos, L. Tolos, "**Equation of state of hot hyperonic neutron star core**", *Proceedings of Science, FAIRness2022* (2023), 027.
- **H. Kochankovski**, À. Ramos, L. Tolos, S. Blacker, A. Bauswein, "**Hyperons in neutron star mergers**", *Proceedings of Science, QNP2024* (2025), 151.
- L. Tolos, **H. Kochankovski**, A. Ramos, G. Lioutas, S. Blacker, A. Bauswein "**Hyperonic degrees of freedom in binary neutron star mergers**", *Submitted*.
- **H. Kochankovski**, À. Ramos, L. Tolos, "**Impact of hyperons on neutron star mergers**", *In preparation*.

Contents

Acknowledgements	i
Resum	iii
Abstract	vii
Glossary	ix
List of publications	xi
1 Introduction	3
1 The birth of neutron stars (NSs)	5
2 The structure of NSs	7
3 Multi-messenger astronomy of cold NSs	11
4 Binary neutron star mergers: a dynamic probe	13
2 Relativistic description of hypernuclear matter at finite temperature	16
1 Relativistic Mean-Field framework	19
1.1 Nucleonic model parameters	25
1.2 Hyperonic model parameters	27
2 Composition and EoS of hypernuclear matter	30
3 Thermal quantities in β –stable NS matter	33
3 Hyperonic effects on neutron star observables	43
1 Mass and radius of cold and hot NSs	44
2 Tidal deformability of cold and hot NSs	49
3 Moment of Inertia of cold and hot NSs	52
4 Strangeness number of cold and hot NSs	55
4 Hyperons signatures in binary neutron star (BNS) mergers	58
1 Thermal behavior of hyperonic equations of state (EoSs)	60

1.1	Brief description of hyperonic EoSs	61
1.2	Density dependence of thermodynamic quantities	64
2	BNS merger simulations	70
2.1	Approach	70
2.2	Maximum density, hyperon fraction, thermal index and temperature evolution in BNS merger simulations	72
3	BNS observables	77
3.1	Dominant frequency peak f_{peak}	78
3.2	Secondary frequency peaks	92
3.3	Mass ejecta	95
3.4	Threshold mass	98
3.5	Further analysis: Effect of the total mass and the asymmetry in BNS mergers	104
5	Summary and Conclusions	110
A	Expansion of the cold nuclear equation of state	115
B	Interplay of hyperon potentials	118
C	Cold star properties of different EoSs and simulation results	125
D	Gravitational wave spectra	130
E	Toy model of nuclear and hypernuclear matter with low thermal index	132
	List Of Figures	135
	Bibliography	143

Chapter 1

Introduction

"Death is the start of a beginning, life is the start of the end, and the cycle continues forever." - Tim I. Gurung

It is unusual for a story to begin with death. But the death of a star is where our journey starts.

The first half of the previous century was one of the most fruitful periods in the history of physics. In the 1920s, while electrons [1] and protons [2] were known, the structure of the nucleus remained mysterious. Rutherford predicted a neutral particle, initially envisioned as a proton-electron pair being part of the nucleus [3], but the initial investigations contradicted this claim [4, 5]. In 1932, Chadwick, building on the work by Bothe [6] and Joliot-Curie [7], definitively discovered the neutron, a neutral particle with a mass similar to the proton [8], resolving the nuclear composition puzzle and earning him the Nobel Prize in 1935.

The discovery of the neutron triggered many different hypotheses. Subjectively speaking, the most interesting ones belonged to the field of astrophysics. The eccentric¹ Swiss physicist Zwicky and his colleague Baade had already predicted that a catastrophic explosion of massive stars could be a source of cosmic rays [9, 10, 11]. Motivated by the discovery of the neutron, they proposed that massive stars end their lives in stupendous explosions that blow them apart and that their remnants represent collapsed stars made of densely packed neutrons [12, 13, 14]. Those neutrons would be created by crushing together the protons and electrons in the gravitational collapse, leaving an object with densities similar to those found in the nuclei and with dimensions of several kilometers. The main force keeping the star from gravitational collapse would then come from the

¹The word "eccentric" is definitely an understatement here. On one occasion, Fritz Zwicky said the following about his fellow colleagues: *"Astronomers are spherical bastards. No matter how you look at them, they are just bastards."*

pressure of strongly degenerate neutrons. Some other authors quickly supported this hypothesis [15, 16, 17, 18, 19]. The idea of Neutron Stars (NSs) as ashes of massive luminous stars was born.

Although physically well motivated, the possible existence of NSs did not encourage extensive research in this area. One would have to wait almost 30 years for progress to be made. In 1967 Bell and *her* advisor Hewish² using a radio telescope, detected a repeating signal in the form of a pulse that they firstly attributed to the existence of "little green men" [20]. It turned out that their discovery was just marginally less exciting. The "pulsar" object whose signal they detected was actually the first detection of a NS, confirming Pacini's prediction published just a few months earlier that NS are in fact strong sources of electromagnetic radiation [21, 22].

Currently, almost 3200 NSs have been detected [23]; most of them are through radio telescopes and are called pulsars. However, some NSs that are in the early stages of evolution can be radio-quiet objects [24, 25]. We now know that NSs are the densest objects in the Universe. Their exotic properties are unmatched – they can be twice as heavy as our Sun, and at the same time they have radii of around 10-15 km [26, 27, 28, 29, 30, 31, 32]. This makes NSs natural laboratories for studying matter in extreme conditions. Furthermore, they are one of the strongest sources of magnetic fields in the Universe, with magnetic fields that can reach up to 10^{12} T. NSs with the strongest magnetic fields are called magnetars [33, 34, 35] and typically are detected as Anomalous X-ray Pulsars and Soft Gamma Repeaters, depending on the type of radiation that they emit [36, 37, 38, 39, 40, 41]. After their formation, NSs do not have a heating source and cool down in a process that lasts millions of years [42, 43, 44, 45, 46, 47, 48]. However, they may still evolve through the accretion of matter [49, 50, 51] or collisions [52, 53].

Although NSs are generally composed of neutrons, a detailed description of their composition and structure, which are driven by the so-called neutron matter Equation Of State (EoS), is far from a trivial question, and we will try to answer it in this thesis. The simplest models assume that in addition to neutrons, NSs have a small amount of protons and electrons. However, the environment found in the NSs could allow the production of exotic species in dense matter. In fact, this was predicted even before the first detection of NSs. In 1960 and 1961, two Soviet (or better said, Armenian) astrophysicists, Ambartsumyan and Saakyan, wrote seminal papers [54, 55] about the possibility of the appearance of hyperons, baryons with strange content in dense matter. As we shall see, decades later, the question of the composition of NSs is still not answered. We will study the effect of the possible production of hyperons on the properties of dense matter and consequently on all observables that are measured nowadays from NSs as a central topic of the thesis. In particular, we will investigate the properties of hot

²Hewish, and not Bell, was awarded the Nobel Prize in Physics in 1974 " for his decisive role in the discovery of pulsars." Bell was not given the real credit that she deserved.

hypernuclear matter in isolated stars and in dynamical phenomena. On the one hand, we will study the effects induced by hyperons found in hot matter, such as the one found when two NSs from a binary collide. On the other hand, we will see what studying the properties of NSs can teach us about the interaction between the particles in extreme conditions, unattainable on Earth.

We begin the thesis with a chapter that provides an introduction to the topic. The process of creation of NSs is briefly reviewed, followed by a description of the structure of a NS. We also discuss in which ways one could extract observables from NSs. A short introduction to Binary Neutron Star (BNS) mergers is provided too. In Chapter 2, the theoretical basis for studying dense matter is outlined. The focus will be given on developing an EoS in a Relativistic Mean Field (RMF) framework for baryonic matter at finite temperature. We will explain how the information of finite nuclei, hypernuclear, and nuclear matter constrains the model parameters. In addition, we will also explore the uncertainties that are present in our model parameters related to the hyperons due to the lack of hypernuclear data. The thermodynamical quantities, such as the pressure, energy, and entropy, will be inspected, with a special emphasis on the behavior of the thermal index when hyperons are present in dense matter. In Chapter 3 we make the connection between the microscopic EoS that is governed by the underlying interaction and the global properties of the stars. Specifically, we calculate the mass-radius, tidal deformability, and moment of inertia of NSs with our developed model as a way to test the predictions with future observations. Both isolated cold stars and young, hot stars are studied. Finally, in Chapter 4, we will move away from a one-model description of matter, and we will first try to find systematic effects that hyperons induce at the EoS level, irrespective of the theoretical framework or parameterization used. The main emphasis is given to the effects found at finite temperature, notably the effect on the thermal pressure, thermal index, and specific heat in matter. Then, we will see how those effects influence the observables and quantities of BNS mergers by studying the results from an extensive simulation campaign. The hyperon imprints on the gravitational wave spectra, the ejecta mass, and the threshold mass are presented. In the end, a summary is also provided.

1 The birth of neutron stars (NSs)

The main-sequence stars, like our Sun, sustain themselves by converting hydrogen into helium ³ through nuclear fusion within their dense cores [57, 58, 59]. Stars significantly more massive than our Sun, specifically those exceeding 8 solar masses ($M > 8M_{\odot}$), embark on a more complex evolutionary path [60, 61]. These massive stars undergo a

³I want to highlight the PhD thesis of *Cecilia Payne-Gaposchkin* [56] where, for the first time, it was predicted that the stars are mostly made of hydrogen and helium. Her thesis represents the birth of modern astrophysics.

series of successive nuclear fusion stages, burning heavier elements in their cores. This progression moves from hydrogen to helium, then carbon, oxygen, neon, and finally silicon (Figure 1.1). The outward radiation pressure produced by these fusion reactions counteracts the inward pull of gravity, leading to the formation of a layered, onion-like structure. Ultimately, this process culminates with the creation of an iron core. Iron represents the final product of nuclear fusion in these stars, as the synthesis of heavier elements requires an input of energy rather than releasing it.

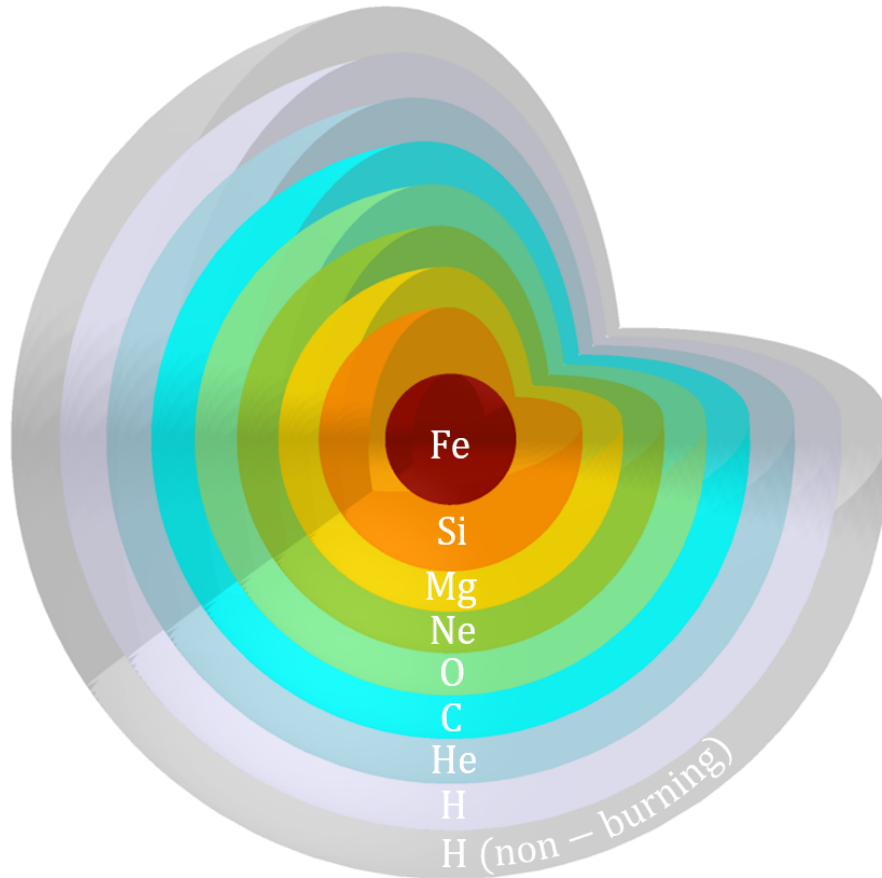


Figure 1.1: Scheme of a massive star's structure at the end of its life as a star on the main sequence. The layers' thickness is not drawn to scale.

Once the mass of the iron core surpasses the Chandrasekhar limit (approximately $M > 1.4M_{\odot}$) [62, 63, 64], the balance between gravity and outward pressure (radiation and electron degeneracy pressure) breaks down. This initiates an implosion that lasts about one second. During the collapse of the core, the rapidly increasing density triggers neutronization, wherein electrons are captured by protons to produce neutrons.

This process occurs when the density of matter exceeds the normal saturation density.⁴ The implosion is followed by a "bounce," which is produced by the matter becoming incompressible. The bounce generates a powerful shock wave, resulting in a type II supernova explosion. As a consequence, a significant amount of matter becomes ejected material that is separated from the remnant. Within the first second, the star's radius dramatically shrinks from roughly 150 km to about 20 km, accompanied by significant mass accretion of the outer layer infalling material for up to 0.5 seconds [65, 66, 67, 68, 69, 70]. The star's subsequent evolution depends on its initial mass. A star with initial mass $M < 25M_{\odot}$ evolves into a Proto-Neutron Star (PNS), the initial phase of its existence as a NS [71, 72]. The main star's phase of life as a PNS is called the Kelvin-Helmholtz epoch, during which significant changes in the structure and the composition of the star occur. During this process, the remnant changes from a hot (the temperature can exceed 10^{11} K) and lepton-rich PNS to a cold and deleptonized NS. In the initial phase, known as the deleptonization stage, excess trapped electron neutrinos diffuse outward from the star's core. The diffusion significantly heats the core of the star while, at the same time, decreasing the proton and electron fraction. This phase is followed by the second stage, known as the cooling stage, in which the star's entropy is slowly decreasing. After the termination of the Kelvin-Helmholtz epoch, around 30 seconds after the birth of the star, most of the deleptonization is completed and the star is cooled down significantly. The PNS settles into a stable configuration and has become a NS. We should note that the star will continue evolving (cool down; it may slow down, and its magnetic field can decay) over millions of years. Eventually, it will become a cold and dense remnant [66, 45, 73]. A summary of the whole evolution of a PNS star is schematically shown in Fig. 1.2.

2 The structure of NSs

The structure of a NS at zero temperature is complex and consists of outer and internal regions (see Fig. (1.3)). The outer region is also known as the atmosphere and represents a plasma layer with a thickness of tens of centimeters. Although the effects of the atmosphere on the hydrodynamical equilibrium of the star are negligible, it is the place where the thermal electromagnetic radiation is formed. By comparing atmospheric observations with the theoretical description of radiation transfer, one may learn about the surface parameters, such as its effective temperature, gravity, and chemical composition. In addition, information about the global properties, such as the mass and radius of the star, can also be extracted [77, 78].

The internal NS regions have a much richer structure than the atmosphere. They can be categorized as core or crustal, depending on whether homogeneous nuclear matter is

⁴approximately a particle density of $\rho_0 = 0.16 \text{ fm}^{-3}$ or a mass density of around $\rho_{M,0} = 2.7 \cdot 10^{14} \text{ g/cm}^3$.

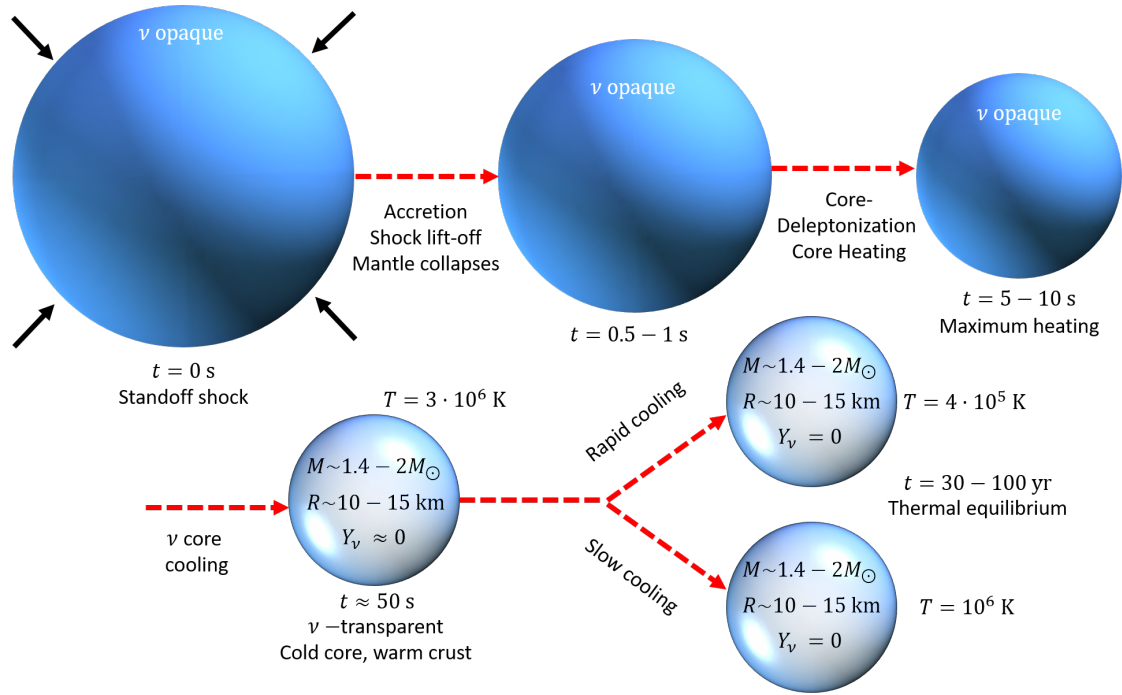


Figure 1.2: Evolution path of a PNS ($M < 25M_{\odot}$) after a successful explosion. It is assumed that the star will not reach the mass for a black hole collapse at any stage. The figure is inspired by similar figures in [74] and in [75] (as a chapter of [76]).

present or not. The thinner crustal part has a thickness of more than a kilometer. When the crustal matter is found at zero temperature, it can be subdivided into an outer and an inner crust. The outer crust is composed of a crystal lattice of nuclei within a sea of electrons. The nuclear species composing the outer crust change toward neutron-rich ones as one moves toward the center of the star [79, 80]. As the density increases, at one point the nuclei become so neutron-rich that a further increase of the density leads to dripping out of the newly formed neutrons. This formation of a neutron gas marks the beginning of the inner crust.

The inner crust is composed of neutron-rich nuclei within a sea of neutrons and electrons. The concentration of free neutrons increases as the density increases [81, 82], and so does the deformation of the nuclei. A distinct state of matter in which nucleons are arranged in a variety of non-spherical clustered structures called the "Pasta Phase" could also occur [83, 84, 85, 86]. At a certain density, around $\rho_B \approx 0.5\rho_0$, the nuclei stop existing, and only homogeneous matter is present. This point represents the beginning of the star's core.

The core of a NS is the central part where the density may even exceed $10\rho_0$. Similarly to the crust, it can be subdivided into two parts. The first one is known as the outer core, and

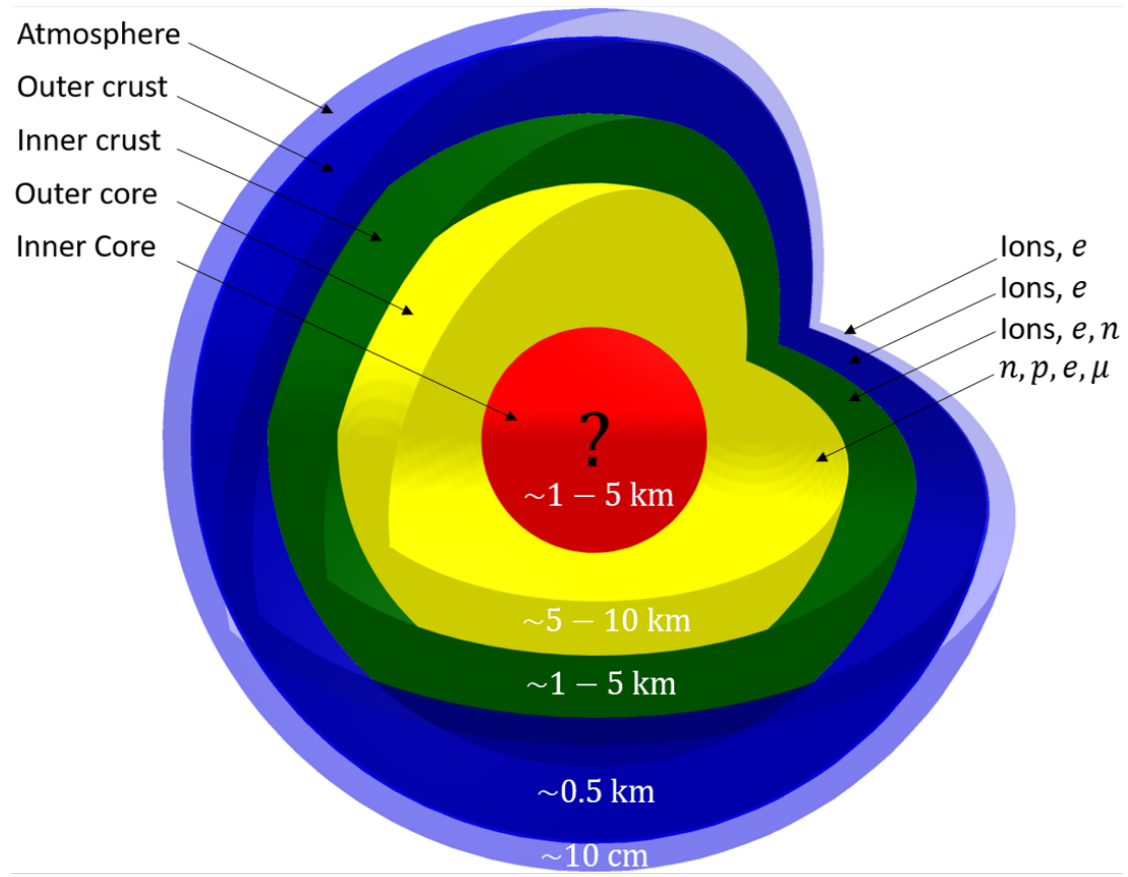


Figure 1.3: Scheme of the structure of a massive cold NS. Please note that the exact thickness of each layer depends on the mass of the star and the interaction assumed between the species (see the discussions in Chapter 3 and Chapter 4.)

its homogeneous matter is mainly composed of neutrons along with protons, electrons, and muons. The latter have been created by the decay of strongly energetic electrons. The leptons can be considered as free Fermi gases. The neutrons and protons, on the other hand, because of the nuclear interaction, form strongly interacting Fermi liquids. Because of the attractive nature of the nuclear force in the medium at long range and the extreme degeneracy, the neutron and proton gases are in a superfluid and superconducting state. They can flow without friction as the bulk and shear viscosity processes are hindered. [45, 87] The beta-equilibrium reaction, along with charge neutrality and baryon number conservation, determines the composition of matter. For the lightest NSs, the outer core stretches all the way up to the center of the star. However, for more massive NSs, the density of matter exceeds $\rho_B \approx 2\rho_0$, which marks the beginning of the inner core. Since it is difficult to reproduce the conditions from the inner core in terrestrial experiments,

the composition of the inner core is still subject to large uncertainties. Different models predict different constituents in the ultra-dense matter [29, 28, 31]. Whereas some of the models assume that the matter is composed of nucleons and leptons only, in many works in the past it has been shown that the production of exotic species could be energetically favorable. In particular, a special emphasis has been put on the possible appearance of deconfined quark matter [88, 89, 90, 91, 92, 93, 94], pion condensation [95, 96, 97, 98, 99, 100, 101], and kaon condensation [102, 103, 104, 105, 106, 107, 108]. Nevertheless, the most promising candidates to appear in the dense matter of the inner core are hyperons. As their potential appearance is the focus of this thesis, we will thoroughly motivate the idea.

Hyperons are baryons similar to nucleons, but they contain at least one strange quark. They obey Fermi-Dirac statistics and can also interact via the strong nuclear force. The lightest hyperons, with a spin-parity $J^P = 1/2^+$, together with the nucleons, form the known baryon octet (see Fig. 1.4 and Table 1.1). They are unstable in terrestrial conditions. The lightest one, the Λ hyperon, has a lifetime of around $2.6 \cdot 10^{-10}$ s and usually decays to a nucleon and a pion:

$$\begin{aligned}\Lambda &\longrightarrow p + \pi^- \\ \Lambda &\longrightarrow n + \pi^0\end{aligned}$$

However, in the dense matter of NSs, these decays may be forbidden. To understand this, let us assume that a NS core is only made of non-interacting degenerate neutrons. As the baryon density increases going toward the center of the star, the neutrons occupy higher energy levels. If the density is high enough, it may happen that the energy of the neutron at the top of the Fermi sea surpasses the rest mass energy of the Λ hyperon. Then, the creation of Λ hyperons is possible through reactions such as

$$n + n \longrightarrow n + \Lambda$$

The created Λ hyperon cannot decay to a neutron because all neutron energy levels below the Λ energy are filled. In equilibrium, the reactions of creation and destruction of Λ hyperons are balanced, which means that the chemical potentials of both species need to be equal, $\mu_n = \mu_\Lambda$. Thus, Λ hyperons can be stable in dense matter. By analogy, other hyperons can be created in matter as well.

Although naive, this example gives an intuitive reason why hyperons can be present in the ultra-dense matter of NSs. Of course, a more detailed explanation of the appearance of hyperons in matter will be given in Chapter 2, where the interaction between the species will be taken into account. We will also see that temperature has a significant influence on the portion of hyperons present in matter.

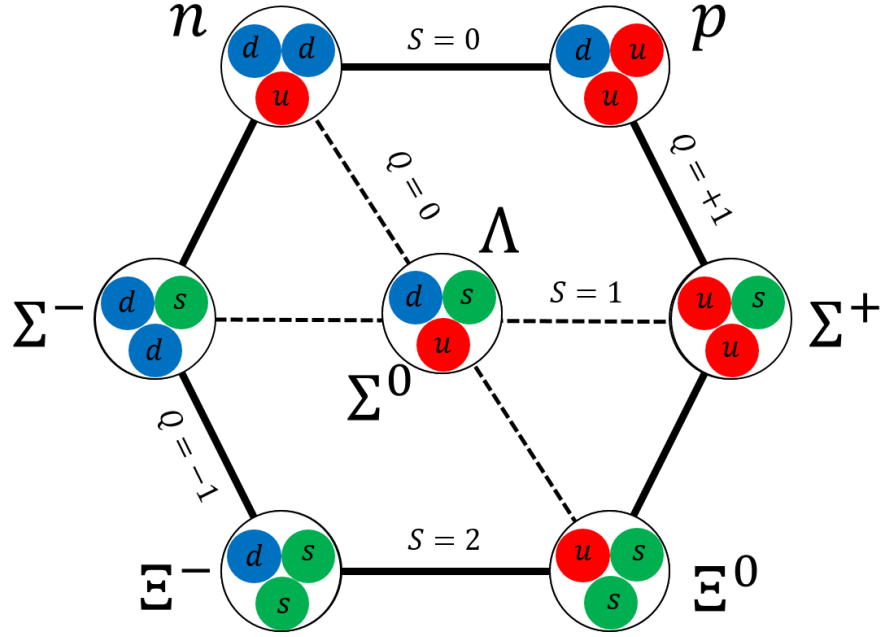


Figure 1.4: Baryon octet. With S the strangeness of the baryons is labeled, while Q represents the baryon electric charge.

Particle	Charge (e)	Mass (MeV)	Quark Content	Mean Lifetime (s)	Most Probable Decay Mode (%)
Proton (p)	+1	938.27	uud	Stable	Stable
Neutron (n)	0	939.57	udd	880.2	$p + e^- + \bar{\nu}_e$
Lambda (Λ^0)	0	1115.68	uds	$2.63 \cdot 10^{-10}$	$p + \pi^-$ (63.9%)
Sigma (Σ^+)	+1	1189.37	uus	$8.02 \cdot 10^{-11}$	$p + \pi^0$ (51.57%)
Sigma (Σ^0)	0	1192.64	uds	$7.4 \cdot 10^{-20}$	$\Lambda^0 + \gamma$ (100%)
Sigma (Σ^-)	-1	1197.45	dds	$1.48 \cdot 10^{-10}$	$n + \pi^-$ (99.8%)
Xi (Ξ^0)	0	1314.86	uss	$2.90 \cdot 10^{-10}$	$\Lambda^0 + \pi^0$ (99.5%)
Xi (Ξ^-)	-1	1321.71	dss	$1.64 \cdot 10^{-10}$	$\Lambda^0 + \pi^-$ (99.9%)

Table 1.1: Properties of the Baryon octet. For more details, see [109, 110, 111].

3 Multi-messenger astronomy of cold NSs

Perhaps the most important properties of any NS are its mass and radius. As we will see in Chapter 3, the mass-radius relation of NSs is closely related to the exact composition

and the interaction between the species. Hence, simultaneous precise measurements of these two quantities can help to constrain the hadronic models at high densities.

The most precise measurements of the mass of NSs that we have nowadays are performed with techniques that utilize the Shapiro delay. The Shapiro delay (or the gravitational time delay effect) is a General Relativity (GR) phenomenon where the travel time of light is slightly increased as it passes through the curved spacetime near a massive object. The effect is particularly pronounced and measurable in highly inclined (nearly edge-on) binary millisecond radio pulsar systems. In such systems, as the radio pulses emitted from the pulsar pass through the gravitational field of its companion star, they experience a measurable time delay. The delay depends directly on the mass of the companion. It encodes the deviations from the Keplerian motion of the objects in the form of post-Keplerian parameters (e.g., Shapiro shape and Shapiro range). By inferring two of these parameters, it is possible to precisely calculate both the NS's and the companion's masses. The measurements of the massive pulsar masses known today have been performed with this technique [112, 113, 114, 115, 116, 117].

The primary method for inferring a NS radius involves analyzing its thermal surface emission spectrum. This technique is particularly effective for isolated NSs, as their observed spectra are not heavily contaminated by magnetospheric emission. Through atmospheric modeling, the effective temperature of the NS can be determined. Subsequently, if the distance to the star is accurately known, using the Stefan-Boltzmann law, the radius can be inferred. However, modeling the thermal emission from NS atmospheres is rather challenging due to the extreme physical conditions. Furthermore, any uncertainties in the distance measurement will propagate directly into the inferred radius, making the distance a significant source of error if it is not well-constrained [118, 119].

Especially useful are the techniques that can constrain the compactness of the star. The spectra captured during the most intense phases of the strong bursts of the X-ray bursters [120], can be utilized to estimate simultaneously the mass and the radius of NSs. The method involves fitting the observed X-ray spectra with theoretical models of hot NS atmosphere. This process constrains several parameters, such as surface gravity and surface gravitational redshift. From the determined values of these two quantities, the star's compactness can then be derived [119, 121, 122].

Another very promising technique has been widely used in the past decade, and it involves X-ray emission from "hot spots" on the surface of millisecond pulsars. Because of the high rotational speed, the X-ray emission spectra from the non-uniform thermal spots of these NSs have a modulated pulse profile. The profile varies with the rotation phase of the NS. Taking into account the GR effects that influence the observed profiles, one can infer the compactness of the NS [123, 124, 125]. The Neutron Star Interior Composition

Explorer (NICER) X-ray telescope has already used this technique to constrain the mass and radius of several NSs [126, 127, 128, 129, 130, 131, 132, 133].

In the future, major advancement in detecting galactic radio pulsars can be expected with the new radio telescope called Square Kilometer Array (SKA) [134]. SKA, with its enhanced sensitivity and improved timing precision, surpassing current telescopes, will enable the discovery of many more radio pulsars in binary systems. This, in turn, will ultimately allow for more precise determination of measured NS masses by at least a factor of ten. In addition, in the following decade, from the spin-orbit coupling of binary systems, the moment of inertia of binary systems will be determined with 10% precision [135, 136, 137, 138, 139, 140]. A measurement of the moment of inertia, coupled with a precise measurement of the mass of NSs, will allow for determination of the radius of some NSs with a precision of 5% [138].

For a more extensive discussion of the topic, we refer to the reviews [141, 142] and the references therein. In Chapter 3 we will see how one can compute different NS observables knowing the EoS of dense matter, and in turn how the measurements that we have nowadays constrain the hadronic models.

4 Binary neutron star mergers: a dynamic probe

Most of the NSs in the Universe are isolated. However, a few percent of them are part of binary systems with a companion being a main-sequence star, red giant, white dwarf, etc. Most interesting are systems where the companion is another NS or a black hole. We will focus on the BNS system. In these cases, the distance between the companions is reduced due to the emission of Gravitational Waves (GWs) [143, 144, 145]. The gradual inspiral takes place as orbital energy and angular momentum are carried away by GWs. The timescale of the inspiral is inversely proportional to the fourth power of the orbital frequency, leading to a dramatic increase in the rate as the stars approach coalescence. The emitted power, on the other hand, is inversely proportional to the fifth power of the semi-major axis. Hence, the GW strain, which represents the fractional change in distance caused by a passing GW, exhibits an increase in amplitude and frequency as the NSs spiral inward. For those reasons, the signal is also known as a "chirp" signal. The GW strain reaches its maximum amplitude during the final moments of the merger, just before the two NSs coalesce. This peak strain corresponds to the point of maximum gravitational acceleration and the most rapid change in the quadrupole moment of the system. Fig. 1.5 shows a schematic representation of the different stages of the merger.

Today, GW emitted during the late moments of the inspiral phase can be detected with the help of GW detectors, such as The Laser Interferometer Gravitational-Wave Observatory (LIGO), Virgo and Kamioka Gravitational Wave Detector (KAGRA). The

precise waveform of the GW signal encodes information about the masses, spins, and tidal deformabilities of the merging NSs, which, when interpreted, describes the state of matter in this environment [146, 147]. GW observations, therefore, offer a unique probe of fundamental physics. In Chapter 3, the connection between the NS structure and the tidal deformability will be explained.

At the end of the life of the binary, the two NSs will collide. This violent astrophysical phenomenon is one of the most fascinating throughout the Universe. With the future GW detectors like Einstein Telescope (ET) and Cosmic Explorer (CE), we will be able to detect the GWs from the post-merger phase as well. Those signals encode valuable information about the state of the hot, ultra-dense matter. The post-merger frequency features will significantly constrain the finite-temperature EoS of dense matter.

Furthermore, BNS mergers are the main environment where, along with GWs, Short Gamma-Ray Bursts (sGRBs) are produced [148, 149]. Powered by the accretion of matter onto a newly formed black hole or a rapidly rotating NS, these sGRBs represent highly collimated jets of relativistic particles. This was confirmed with the detection of the GW170817 merger event, where, alongside GW emission, the sGRB 170817A burst was also detected [150]. This multimessenger observation provided definitive evidence that mergers are progenitors of sGRBs. In the future, the time delay between the GW signal and the GRB emission could provide insights on the jet launching mechanisms of GRBs [151].

In addition, BNS mergers are the origins of another type of radiation known as kilonova radiation [152, 153]. During the merger of two NSs, an extremely high flux of neutrons could be created, and it is possible for the so-called r-process to occur. The r-process in the mass ejecta is responsible for the creation of the neutron-richest isotopes in nature [154, 155]. The corresponding electromagnetic transient powered by the radioactive decay represents the kilonova radiation. The spectral evolution of the kilonova provides information about the composition and velocity of the ejecta, allowing for an estimation to be done on the amount of heavy elements produced.

As we can see, the BNS merger events are an invaluable tool for studying the matter in extreme conditions. In Chapter 4 we will explore how one can simulate these events and can use the information extracted from actual observations to constrain dense-matter models.

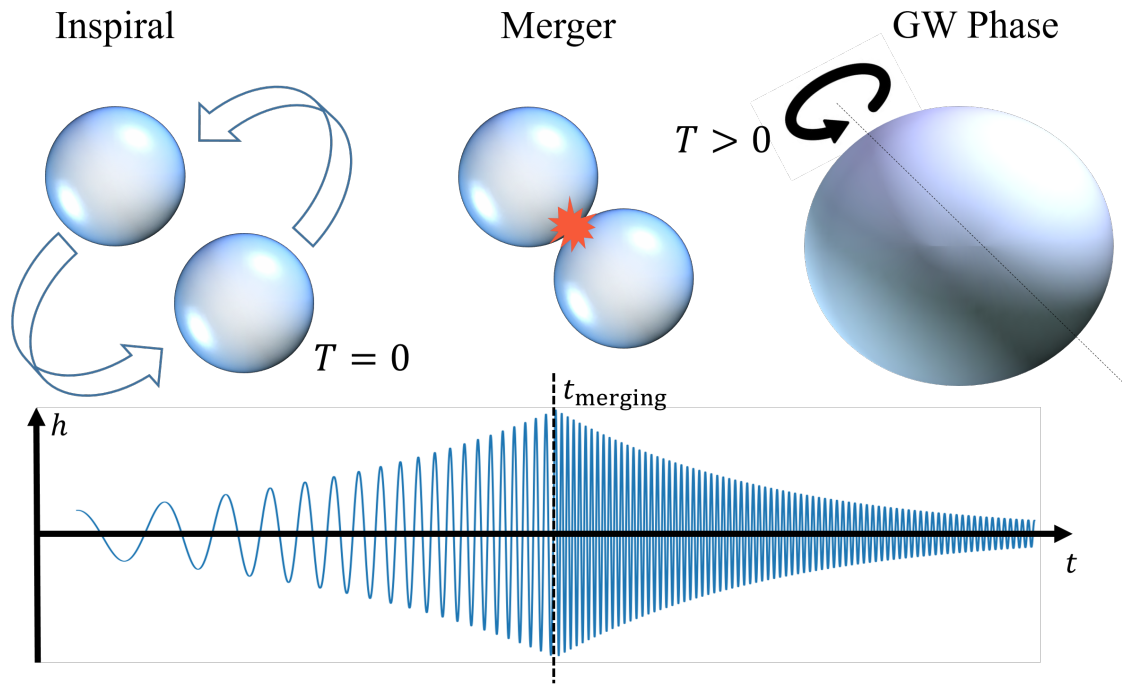


Figure 1.5: Scheme of a BNS merging event in which a hyper-massive NS is formed. Three different phases—inspiral, merger, and the GW phase of the process are marked. The strain of the GWs as a function of time is also shown (amplitude and oscillations not to scale).

Chapter 2

Relativistic description of hypernuclear matter at finite temperature

NSs offer a unique possibility for studying dense matter under extreme conditions. To model a NS, from a theoretical point of view, one needs to know the EoS of NS matter. The EoS is then used as an essential input in the modeling of different NS processes, from their birth to their old age and in NS merger processes. When the term “Equation of State” is used, it refers to the different relations between thermodynamic state equilibrium variables (energy density (ϵ), pressure (P), chemical potentials (μ_i)...). However, because the different thermodynamic quantities are related with thermodynamic relations, one only needs to know how the proper thermodynamic potential changes with the conditions in the system. If the NS matter is cold and in β -equilibrium, as in the evolved and isolated NSs, then

$$\epsilon = \epsilon(\rho_B), \quad (2.1)$$

is the proper thermodynamic potential that describes the system. We see that the cold NS EoS is one-dimensional (1D), as the energy density only depends on the baryon density (ρ_B)¹. However, in general, the NS matter is neither cold nor in β -equilibrium. Hence, we need to find how the free energy density (f) depends on the matter conditions, e.g., how f changes as a function of the density, composition ($[Y]$) and temperature (T) of matter,

$$f = f(\rho_B, [Y], T), \quad (2.2)$$

to fully determine the EoS of the system. The free energy has a minimum value in the ground state of the system for the given set of thermodynamic variables [156], and represents the proper thermodynamic potential.

¹If it is not otherwise specified, by the term “density” we are referring to *number density*.

From the first prediction of the properties of cold NSs, it had been clear that the interaction between the nucleons, or baryons in general, plays a crucial role in the EoS. The naive model of Oppenheimer and Volkoff [18], where the matter in the star was assumed to form a non-interacting Fermi gas, predicted that the most massive NSs would have masses $M < 0.7 M_{\odot}$. This prediction created a paradox, as it contradicted the evolutionary paths of stars that would be massive enough to form NSs in the first place. Hence, since the pioneering work of Cameron [157], the interaction between the baryons has been acknowledged to play a crucial role in any realistic model that describes NS matter.

In principle, the EoS of strongly interacting matter composed of nucleons and hyperons (hypernuclear matter) could be obtained from Quantum Chromodynamics (QCD), as QCD is the fundamental theory of the strong interaction. While theoretically it dictates the baryon-baryon interactions, perturbative QCD struggles to predict the EoS at low energies due to its increasing coupling strength [158, 159]. Consequently, quarks are confined within the baryons, and baryons themselves can be used as relevant building blocks for describing these interactions at lower energies.

In the literature, two different types of approaches have been used to obtain the baryon interactions at zero temperature and consequently to obtain the EoS of cold hypernuclear matter.

The first ones are the *Ab initio* many-body methods that use 'realistic' two- and three-baryon interactions that are fitted to experimental data as their starting point. The realistic forces are usually based on the meson-exchange theory [160, 161, 162, 163, 164, 165, 166] or are built from chiral perturbation theory [167, 168]. The many-body problem is then solved using different techniques. One of the most widely known methods is the Brueckner-Hartree-Fock (BHF) method where an effective interaction, called the G -matrix, is built by summing the multiple scattering processes between two baryons in the medium. The method involves solving the Bethe-Goldstone equation to determine the energy of the system [169, 170, 171, 172, 173, 174, 175, 176, 177, 178]. The Dirac BHF method, on the other hand, extends the BHF method to incorporate relativistic effects, which are essential for accurately describing dense baryon matter [179, 180, 181, 182]. Other notable many-body methods used to obtain the EoS of cold hypernuclear matter are the low-momentum potential V_{low} approach, where the baryon interactions are "softened" by integrating out high-momentum components [183], the Monte Carlo method, where statistical sampling is used for an approximate solution to the hypernuclear problem [184, 185], and the lowest-order constrained variational method that minimizes the system's overall energy functional while adhering to a normalization constraint [186].

The second approach, known as the phenomenological approach, consists of models where the inputs are density-dependent effective interactions that contain a certain number of parameters. Those parameters are adjusted so that the resulting EoS is able to reproduce

experimental data for nuclei all over the nuclear chart as well as nuclear matter properties. Several different frameworks have been developed. The most widely used are the RMF models that we will employ in this work. Starting from Glendening [187, 188, 189, 190], many authors have investigated the properties of hypernuclear matter and cold NSs within this framework [191, 192, 193, 194, 195, 196, 197, 198, 198, 199, 200, 201, 202, 203, 204, 205, 206, 207, 208, 209, 210, 211, 212, 213, 214, 215, 216, 217, 218, 219, 220, 221].

Another popular framework is the non-relativistic Skyrme model. Originally developed by Skyrme in the late 1950s, it serves as an effective interaction for nuclear Hartree-Fock calculations [222, 223]. Later, this approach was also extended to study hypernuclear matter [224, 225]. For infinite hypernuclear matter, the model yields analytical expressions, which is a major advantage. However, at higher densities the relativistic effects start playing a significant role, so the reliability of the Skyrme models decreases.

We refer the interested reader to read about the different frameworks used in the literature in the context of hyperons in cold NS in the following reviews [226, 227, 228, 229, 230, 231, 232, 233]. Note that some models that have been used extensively for the pure nuclear matter problem, e.g. Self-Consistent Green's functions [234, 235, 236] or Gogny forces [237, 238, 239, 240] have not been used to study hypernuclear matter [241, 242].

While the cold EoS is suitable to describe the matter in evolved and isolated NSs, this is not the case for the matter found in violent astronomical phenomena or during the early evolution of the star. In those conditions the temperature can even exceed 100 MeV, so a finite-temperature EoS is necessary [243, 244, 245]. There was significant progress in obtaining the finite-temperature hypernuclear EoS in recent years [246, 247, 248, 249, 250, 251, 252, 253, 254, 255, 256, 257, 258], specifically in phenomenological frameworks. However, there is still a lack of detailed studies on the behavior of thermodynamic quantities in hypernuclear matter at finite temperature. Furthermore, the large uncertainties in hyperonic effective interactions, particularly at high densities and large isospin asymmetries, are not well-addressed by current models. Therefore, it is important to develop an EoS for homogeneous matter that covers a broad spectrum of temperatures ($T = 0 - 100$ MeV), charge fractions ($Y_Q = 0.01 - 0.6$), and baryon densities ($\rho_B = \rho_0/2 - 6\rho_0$) that allows for a systematic study of the impact of the uncertainties on the composition and the thermodynamic quantities. We will address these issues in this chapter.

The remainder of the chapter is devoted to explaining the development of the finite-temperature RMF model for hypernuclear matter based on the FSU2H parametrization scheme [259, 260, 261, 262]. The model provides an EoS that is suitable as an input in relativistic simulations. We also quantify the uncertainty that exists in the parameter models, and we relate it with the uncertainty that is present in the effective Hyperon-Nucleon (YN) and Hyperon-Hyperon (YY) interactions.

The chapter is organized in the following way. First, we describe the formalism needed to build a finite-temperature EoS in the RMF framework. Then, we explain how the model parameters can be constrained to the existent nuclear and hypernuclear physics data. We will study the properties of hypernuclear matter in detail, as they are our primary focus in this chapter. Special attention will be given to the uncertainty in the effective interaction and how it affects the predictions for NS matter composition and the EoS. Finally, we investigate the properties of hot NS matter in β -equilibrium by studying the thermal effects of the EoS.

1 Relativistic Mean-Field framework

We discussed in the introduction of this chapter that the EoS of dense matter, both at zero and at finite temperature, can be obtained in different frameworks. The RMF framework has several advantages over other methods. First, it provides a simple framework for explicitly including the meson degrees of freedom. This is important for NS ultra-dense matter, where baryons can overlap significantly and the meson fields are distorted. As a consequence, many-body forces start to play a significant role which can be easily incorporated in a RMF framework. Second, the momenta of the species at high densities are increased, and the non-relativistic approximations are not valid anymore. Third, different baryonic degrees of freedom can be easily incorporated. Finally, the generalization of the method to finite temperature is straightforward.

The RMF model was firstly outlined by Walecka in 1974, [263]² where he described the interaction between nucleons in nuclei and nuclear matter on the basis of the exchange of virtual mesons. In this formalism, the nucleonic field:

$$\Psi = \begin{pmatrix} n \\ p \end{pmatrix}, \quad (2.3)$$

couples with different mesons through various terms in which the coupling constants are unknown parameters. Those constants are then fitted to experimental observables. The mesons and different terms in the Lagrangian are not known "*a priori*". In the original Walecka model, meant to explain nuclear matter properties, the interaction is mediated by two massive mesons: the isoscalar-scalar meson σ and the isoscalar-vector meson ω . The first one couples to the scalar nucleon density, $\bar{\Psi}\Psi$, while the latter to the nucleon current, $\bar{\Psi}\gamma^\mu\Psi$ [265]. Those interactions are described with Lagrangian terms like $\sigma\bar{\Psi}\Psi$ and $\bar{\Psi}\gamma^\mu\Psi\omega_\mu$. A simple model that accounts only for those interaction channels gives rise to an effective nucleon-nucleon potential that exhibits a short-range repulsion and an intermediate-range attraction, the two main features of the nucleon-nucleon interaction.

²Duerr [264] published a similar idea two decades earlier.

However, soon it has become clear that the model is too simple to capture the complexity of the many-body nuclear problem. The successive models implemented new terms. First, the vector isovector meson $\vec{\rho}$ was introduced to properly explain the behavior of asymmetric matter, and the σ – meson self-interactions ($\kappa\sigma^3$ and $\lambda\sigma^4$) were added to achieve realistic compressibility [266]. Bodmer [267] added the ω -meson quartic self-coupling ($\zeta(\omega_\mu\omega^\mu)^2$) to ensure stability and soften the EoS at high densities, which is crucial for NS modeling. Horowitz and Piekarewicz, [268, 269], incorporated the $\omega - \rho$ meson mixing ($\Lambda_{\omega\rho}\vec{\rho}_\mu\vec{\rho}^\mu\omega_\mu\omega^\mu$) to fine-tune the nuclear symmetry energy predicted by the model.

One can extend the previous nucleonic field to a generic baryonic one that comprises all baryons from the baryon octet:

$$\Psi = \begin{pmatrix} n \\ p \\ \Lambda \\ \dots \\ \Xi^0 \end{pmatrix}, \quad (2.4)$$

Glendenning [187] was the first to consider the hyperonic degrees of freedom. Soon it was clear that additional terms were necessary to realistically describe the YY interactions. The isoscalar-vector hidden meson ϕ was introduced to stiffen the EoS when hyperons are present in the matter, and the isoscalar-scalar σ^* hidden meson was added to fine-tune the YY interactions [193]. We note that an alternative formalism where, instead of introducing self-interacting terms, the coupling constants are density dependent has also been studied in detail in the literature; see, for example, Refs. [270, 271, 272, 273, 274, 214].

To obtain the NS EoS at arbitrary conditions, we need a proper description of the strongly interacting baryonic matter. In our model, we assume that hypernuclear matter is made of strongly interacting baryons ($b = n, p, \Lambda, \Sigma, \Xi$) at a given temperature (T), density (ρ_B), and charge fraction (Y_Q). The charge fraction of the system is related to the density of each baryon ρ_b and its charge Q_b as

$$Y_Q = \sum_b \frac{Q_b \rho_b}{\rho_B}. \quad (2.5)$$

The Lagrangian in our model can be written as

$$\begin{aligned}
\mathcal{L} &= \sum_b \mathcal{L}_b + \mathcal{L}_m, \\
\mathcal{L}_b &= \bar{\psi}_b (i\gamma_\mu \partial^\mu - q_b \gamma_\mu A^\mu - m_b + g_{\sigma b} \sigma + g_{\sigma^* b} \sigma^* \\
&\quad - g_{\omega b} \gamma_\mu \omega^\mu - g_{\phi b} \gamma_\mu \phi^\mu - g_{\rho, b} \gamma_\mu \vec{I}_b \cdot \vec{\rho}^\mu) \psi_b, \\
\mathcal{L}_m &= \frac{1}{2} \partial_\mu \sigma \partial^\mu \sigma - \frac{1}{2} m_\sigma^2 \sigma^2 - \frac{\kappa}{3!} (g_{\sigma N} \sigma)^3 - \frac{\lambda}{4!} (g_{\sigma N} \sigma)^4 \\
&\quad + \frac{1}{2} \partial_\mu \sigma^* \partial^\mu \sigma^* - \frac{1}{2} m_{\sigma^*}^2 \sigma^{*2} \\
&\quad - \frac{1}{4} \Omega^{\mu\nu} \Omega_{\mu\nu} + \frac{1}{2} m_\omega^2 \omega_\mu \omega^\mu + \frac{\zeta}{4!} g_{\omega N}^4 (\omega_\mu \omega^\mu)^2 \\
&\quad - \frac{1}{4} \vec{R}^{\mu\nu} \vec{R}_{\mu\nu} + \frac{1}{2} m_\rho^2 \vec{\rho}_\mu \cdot \vec{\rho}^\mu + \Lambda_\omega g_{\rho N}^2 \vec{\rho}_\mu \vec{\rho}^\mu g_{\omega N}^2 \omega_\mu \omega^\mu \\
&\quad - \frac{1}{4} P^{\mu\nu} P_{\mu\nu} + \frac{1}{2} m_\phi^2 \phi_\mu \phi^\mu - \frac{1}{4} F^{\mu\nu} F_{\mu\nu}.
\end{aligned} \tag{2.6}$$

The Lagrangian is comprised of a baryonic part, $\mathcal{L}_b$, and a mesonic part, $\mathcal{L}_m$. The baryonic part involves the components ψ_b of the Dirac field Ψ , the electromagnetic field vector A^μ , the particle masses m_i , the isospin operator $\vec{I}$ and the Dirac matrices γ_μ . The field strength tensors for the vector mesons and electromagnetic fields are $\Omega^{\mu\nu}$, $\vec{R}^{\mu\nu}$, $P^{\mu\nu}$, and $F^{\mu\nu}$. The couplings between the baryons and the mesons are denoted by g_{mb} . Although the electromagnetic potential is present in the Lagrangian, it will be disregarded because our main focus is to describe charge-neutral systems without magnetic fields. For simplicity, we do not include the explicit dependence of the meson and baryon fields on the spacetime coordinate- x .

The concept of an EoS can only be applied if the studied systems have achieved thermodynamic equilibrium. While thermal and mechanical equilibrium are happening almost instantaneously, the same does not always hold for chemical equilibrium. The chemical equilibrium can be assumed only if the system's hydrodynamic evolution is much longer than the time scale of the reactions. The weak interactions between baryons occur on a time scale of 10^{-6} s [104]. This is shorter than the timescale of the hydrodynamic evolution, both in PNS evolution and in the BNS merger simulation. So, the weak equilibrium between the baryons can be assumed to hold in all phenomena of interest. That raises the following relations between the chemical potentials of the species:

$$\begin{aligned}
\mu_{b^0} &= \mu_n, \\
\mu_{b^-} &= 2\mu_n - \mu_p, \\
\mu_{b^+} &= \mu_p.
\end{aligned} \tag{2.7}$$

In the previous equations, μ_b labels the chemical potential of the baryon species b , and b^0 , b^- , b^+ indicate neutral, negatively charged and positively charged baryons, respectively. Consequently, the strangeness number is not a conserved quantity. The only two conserved charges are the total baryon number:

$$\rho_B = \sum_b \rho_b, \tag{2.8}$$

and the charge density:

$$\rho_Q = \sum_b Q_b \rho_b. \tag{2.9}$$

The presence of leptons is necessary to achieve charge neutrality in matter. However, the weak interactions that include the leptons cannot generally be considered to be in equilibrium. For example, the time scale of the electron-capture reaction:

$$p + e^- \longrightarrow n + \nu_e$$

in some conditions exceeds the dynamical time scale of the process. This is the reason why we do not consider them at the Lagrangian level. However, they can be easily added when an equilibrium condition is fulfilled, as they can be treated as non-interacting Fermi gases.

To determine the system's composition and EoS, we must obtain and solve the Euler-Lagrangian equations of motion. Knowing that for any field ϕ the corresponding equation can be found as:

$$\partial_\nu \left(\frac{\partial \mathcal{L}}{\partial (\partial_\nu \phi)} \right) - \frac{\partial \mathcal{L}}{\partial \phi} = 0 \tag{2.10}$$

the mesonic equations of motion read:

$$\begin{aligned}
\partial_\mu \partial^\mu \sigma + m_\sigma^2 \sigma + \frac{\kappa}{2} g_{\sigma N}^3 \sigma^2 + \frac{\lambda}{3!} g_{\sigma N}^4 \sigma^3 &= \sum_b g_{\sigma b} \bar{\psi}_b \psi_b, \\
\partial_\mu \partial^\mu \sigma^* + m_{\sigma^*}^2 \sigma^* &= \sum_b g_{\sigma^* b} \bar{\psi}_b \psi_b, \\
\partial_\nu \Omega^{\mu\nu} + m_\omega^2 \omega^\mu + \frac{\zeta}{3!} g_{\omega N}^4 \omega^\mu \omega_\mu \omega_\nu + 2\Lambda_\omega g_{\rho N}^2 g_{\omega N}^2 \omega^\mu \vec{\rho}^\nu \vec{\rho}_\nu &= \sum_b g_{\omega b} \bar{\psi}_b \gamma^\mu \psi_b, \\
\partial_\nu \vec{R}^{\mu\nu} + m_\rho^2 \vec{\rho}^\mu + 2\Lambda_\omega g_{\rho N}^2 g_{\omega N}^2 \omega^\mu \omega_\mu \vec{\rho}^\mu &= \sum_b g_{\rho b} \bar{\psi}_b \gamma^\mu \vec{I}_b \psi_b, \\
\partial_\nu P^{\mu\nu} + m_\phi^2 \phi &= \sum_b g_{\phi b} \bar{\psi}_b \gamma^\mu \psi_b.
\end{aligned} \tag{2.11}$$

The equations for the scalar mesons are known as Klein-Gordon equations, while the ones for the vector mesons are known as Proca equations [275]. These equations are nonlinear coupled differential equations that are difficult to solve. However, the equations can be simplified significantly. We notice that the source terms in the equations are related with the conserved baryon density. For high-density systems the quantum fluctuations can be neglected and a suitable approximation is to replace the meson fields by classical fields and the sources with their expectation values. Furthermore, since the system is static, the baryon fluxes $\psi_b \gamma^i \psi_b$ are 0, so the non-zero components of the vector mesons vanish. As for the isovector mesons like $\vec{\rho}$, since the matter is considered as isotropic, all component except the third ones are vanishing³. These simplifications are known as RMF approximation:

$$\begin{aligned}
\sigma &\rightarrow \bar{\sigma} = \langle \sigma \rangle; \\
\vec{\rho} &\rightarrow \bar{\rho} = \langle \rho_3^0 \rangle; \\
\omega &\rightarrow \bar{\omega} = \langle \omega^0 \rangle; \\
\phi &\rightarrow \bar{\phi} = \langle \phi^0 \rangle; \\
\sigma^* &\rightarrow \bar{\sigma}^* = \langle \sigma^* \rangle.
\end{aligned} \tag{2.12}$$

On the other hand, the equations of motion for the baryon fields are given by Dirac equations:

$$\left(\gamma^\mu (i\partial_\mu - g_{b\omega} \bar{\omega} - g_{b\rho} \bar{\rho} - g_{b\phi} \bar{\phi}) - (m_b - g_{b\sigma} \bar{\sigma} - g_{b\sigma^*} \bar{\sigma}^*) \right) \psi_b = 0. \tag{2.13}$$

³Especially when dealing with exotic nuclear phenomena or nuclear systems, one might want to relax the assumptions of spatial homogeneity and isotropy so other components of the ρ meson could become relevant [276].

We note that the scalar fields modify the mass of the particle, leading to an effective mass:

$$m_b^* = m_b - g_{b\sigma}\bar{\sigma} - g_{b\sigma^*}\bar{\sigma}^*, \quad (2.14)$$

while the vector fields are affecting the energy, which give rise to modification of the true chemical potential of the species, μ_b , into an effective chemical potential

$$\mu_b^* = \mu_b - g_{b\omega}\bar{\omega} - g_{b\rho}I_{3b}\bar{\rho} - g_{b\phi}\bar{\phi}. \quad (2.15)$$

The baryon fields, even in a mean-field approximation, remain operators. In Eqs. (2.11), the source terms must be evaluated on the ground state since they are the source term origin for the meson fields. At finite temperature, this leads to [277]:

$$\begin{aligned} \bar{\psi}_b \psi_b &\rightarrow \langle \bar{\psi}_b \psi_b \rangle = \rho_s = \frac{\gamma_b}{2\pi^2} \int_0^\infty dk k^2 \frac{m_b^*}{\sqrt{k^2 + m_b^{*2}}} (f_b(k, T) + f_{\bar{b}}(k, T)); \\ \bar{\psi}_b \gamma^\mu \psi_b &\rightarrow \langle \bar{\psi}_b \gamma^0 \psi_b \rangle = \rho_b = \frac{\gamma_b}{2\pi^2} \int_0^\infty dk k^2 (f_b(k, T) - f_{\bar{b}}(k, T)). \end{aligned} \quad (2.16)$$

In the last equations, $\gamma_b = 2$ represents the spin degeneracy of baryons and

$$f_{b(\bar{b})}(k, T) = \left[1 + \exp \left(\frac{\sqrt{k^2 + m_{b(\bar{b})}^{*2}} - \mu_{b(\bar{b})}^*}{T} \right) \right]^{-1} \quad (2.17)$$

represents the Fermi-Dirac distribution for baryon(b)/antibaryon($\bar{b}$) with effective masses $m_b^* = m_{\bar{b}}^*$, and effective chemical potentials $\mu_b^* = -\mu_{\bar{b}}^*$. Finally, the equations of the meson fields can be written as:

$$\begin{aligned} m_\sigma^2 \bar{\sigma} + \frac{\kappa}{2} g_{\sigma N}^3 \bar{\sigma}^2 + \frac{\lambda}{3!} g_{\sigma N}^4 \bar{\sigma}^3 &= \sum_b g_{\sigma b} \rho_b^s, \\ m_{\sigma^*}^2 \bar{\sigma}^* &= \sum_b g_{\sigma^* b} \rho_b^s, \\ m_\omega^2 \bar{\omega} + \frac{\zeta}{3!} g_{\omega N}^4 \bar{\omega}^3 + 2\Lambda_\omega g_{\rho N}^2 g_{\omega N}^2 \bar{\omega} \bar{\rho}^2 &= \sum_b g_{\omega b} \rho_b, \\ m_\rho^2 \bar{\rho} + 2\Lambda_\omega g_{\rho N}^2 g_{\omega N}^2 \bar{\omega}^2 \bar{\rho} &= \sum_b g_{\rho b} I_{3b} \rho_b, \\ m_\phi^2 \bar{\phi} &= \sum_b g_{\phi b} \rho_b. \end{aligned} \quad (2.18)$$

Eqs. (2.18) together with Eqs. (2.7), (2.8), (2.9), (2.15), (2.16) form a set of coupled non-linear integral equations that fully determine the composition of matter. The solution

of the system yields the fields of the mesons, the chemical potentials of the baryons, and the individual densities of each species. The stress-energy momentum tensor:

$$T_{\mu\nu} = \frac{\partial \mathcal{L}}{\partial(\partial_\mu \phi_\alpha)} \partial^\mu \phi_\alpha - \eta_{\mu\nu} \mathcal{L}, \quad (2.19)$$

directly allows for the calculation of various thermodynamic variables. First, one can calculate the energy density of the system:

$$\begin{aligned} \epsilon_B &= \langle T_{00} \rangle \\ &= \frac{1}{2\pi^2} \sum_b \gamma_b \int_0^\infty dk k^2 \sqrt{k^2 + m_b^{*2}} (f_b(k, T) - f_{\bar{b}}(k, T)) \\ &\quad + \frac{1}{2} (m_\omega^2 \bar{\omega}^2 + m_\rho^2 \bar{\rho}^2 + m_\phi^2 \bar{\phi}^2 + m_\sigma^2 \bar{\sigma}^2 + m_{\sigma^*}^2 \bar{\sigma}^{*2}) \\ &\quad + \frac{\kappa}{3!} (g_{\sigma N} \bar{\sigma})^3 + \frac{\lambda}{4!} (g_{\sigma N} \bar{\sigma})^4 + \frac{\zeta}{8} (g_{\omega N} \bar{\omega})^4 + 3\Lambda_\omega (g_{\rho N} g_{\omega N} \bar{\rho} \bar{\omega})^2. \end{aligned} \quad (2.20)$$

Similarly, the pressure can be calculated as:

$$\begin{aligned} P_B &= \frac{1}{3} \langle T_{jj} \rangle \\ &= \frac{1}{6\pi^2} \sum_b \gamma_b \int_0^\infty dk \frac{k^4}{\sqrt{k^2 + m_b^{*2}}} (f_b(k, T) - f_{\bar{b}}(k, T)) \\ &\quad + \frac{1}{2} (m_\omega^2 \bar{\omega}^2 + m_\rho^2 \bar{\rho}^2 + m_\phi^2 \bar{\phi}^2 - m_\sigma^2 \bar{\sigma}^2 - m_{\sigma^*}^2 \bar{\sigma}^{*2}) \\ &\quad - \frac{\kappa}{3!} (g_{\sigma N} \bar{\sigma})^3 - \frac{\lambda}{4!} (g_{\sigma N} \bar{\sigma})^4 + \frac{1}{24} \zeta (g_{\omega N} \bar{\omega})^4 + \Lambda_\omega (g_{\rho N} g_{\omega N} \bar{\rho} \bar{\omega})^2. \end{aligned} \quad (2.21)$$

Finally, the entropy density, s_B , and the free energy density of the system, f_B , can be obtained using the well-known thermodynamic relations [278]:

$$s_B = \frac{1}{T} \left(\epsilon_B + P_B - \sum_i \mu_i \rho_i \right); \quad (2.22)$$

$$f_B = \epsilon_B - s_B T. \quad (2.23)$$

1.1 Nucleonic model parameters

The described formalism, for a given set of parameters, yields an EoS of dense matter at finite temperature. However, the obtained EoS needs to be consistent in reproducing

the quantities and the constraints derived from nuclear experiments. At the same time, the EoS needs to be compatible with the astrophysical observables that come from the multi-messenger observations. The model that we present here, FSU2H*, shares the same nucleonic core as its predecessor, FSU2H. The details about the choice of the parameters in the FSU2H model has been previously discussed in Refs. [259, 260]), and we will just summarize it here.

The FSU2H model was built on the grounds of the FSU2 model [279] that was able to predict heavy cold NSs ($M > 2M_{\odot}$) with large radii, not compatible with recent observations [129, 130]. In an attempt to address this issue, an increase of the Λ_{ω} coupling was proposed, which in turn made the resulting EoS softer in the low- to intermediate-density region ($\rho_0 - 2.0\rho_0$). To counteract the overall softening and in order to be able to predict NSs with high masses, consistent with astrophysical observations, the ζ coupling of the model had to be decreased. This resulted in a stiffer EoS at higher densities. Of course, all other coupling constants needed to be refitted as the EoS had to be compatible with the experimental constraints including nuclei data. The final values of the coupling constants are given in Table 2.2.

The resulting EoS was fixed to reproduce the value of the saturation density in Symmetric Nuclear Matter (SNM) $\rho_0 = 0.1505 \text{ fm}^{-3}$ and the value of the binding energy at saturation density of $E_0 = -16.28 \text{ MeV}$. The value of the saturation density is in agreement with the results obtained from the PREX collaboration [280], as well as with some older theoretical constraints [281, 282]. The value of the binding energy at saturation density in SNM is close to the value obtained by [283] in an analysis of the non-relativistic semi-empirical four-parameter mass formula with coefficients from the shell-corrected Thomas–Fermi model.

For a highly isospin asymmetric matter as the one found in NSs, it is crucial that the EoS is consistent with the constraints obtained for the symmetry energy $E_{\text{sym},0}$, its slope L and the pressure in Pure Neutron Matter (PNM) at saturation density P_{PNM} . Indeed, the value that our model predicts for $E_{\text{sym},0} = 30.5 \text{ MeV}$ is compatible with the results presented in [284]. The same holds for the value of $P_{\text{PNM}}(\rho_0) = 2.3 \text{ MeV/fm}^3$ [285] and for $L = 44.5 \text{ MeV}$ [284, 286, 287]. However, the value of L is substantially lower than the value extracted from the PREX-II analysis [288]. Moreover, in many analysis in the past decades, the interplay between the slope of the symmetry energy and the symmetry energy has been extensively studied, and several strong constraints have been obtained (see Fig. 1 in [259]), that the FSU2H model is able to fulfill. As one can see, the model is consistent with the experimental nuclear data.

For the proper definition of the parameters of nuclear matter and their relation with the EoS, we refer the reader to Appendix A.

1.2 Hyperonic model parameters

The wealth of experimental data and *ab initio* calculations characterizing nuclear matter at densities around saturation density, together with the constraints from astrophysical observations, allowed us to constrain the nucleonic coupling constants of the FSU2H* model with high accuracy. However, a significant knowledge gap persists regarding hypernuclei and hypernuclear matter. Consequently, the couplings of hyperons to vector mesons ω , ρ , and ϕ are obtained from the nucleonic coupling constants assuming SU(3) symmetry, supplemented by the vector dominance model and the assumption of ideal mixing for the physical ω and ϕ mesons [193, 204, 209, 272].

Within this framework, the free parameters of the model are reduced to the couplings of the σ meson with all baryonic species and the coupling of the σ^* meson specifically with the Λ hyperon. Hyperon- σ couplings are determined by fitting the different hyperon species potentials

$$U_b = -g_{\sigma b}\bar{\sigma} - g_{\sigma^* b}\bar{\sigma}^* + g_{\omega b}\bar{\omega} + g_{\rho b}\bar{\rho}I_{3b} + g_{\phi b}\bar{\phi}, \quad (2.24)$$

in SNM, while the Λ - σ^* coupling is obtained from the $\Lambda\Lambda$ bond energy

$$\Delta B_{\Lambda\Lambda}(^A_{\Lambda\Lambda}Z) = B_{\Lambda\Lambda}(^A_{\Lambda\Lambda}Z) - 2B_{\Lambda}(^A_{\Lambda}Z), \quad (2.25)$$

in double- Λ hypernuclei. The $g_{b\sigma}$ couplings of the FSU2H* model were obtained in [289] by fitting the potentials to the widely accepted central values of hypernuclear potentials. The Λ hyperon potential at nuclear saturation density in SNM is particularly well established. In accordance with [290], we employed a value of $U_{\Lambda} = -28$ MeV, which was derived by examining data from the level spectra of Λ hypernuclei generated at emulsion and bubble chambers through reactions (π^+, K^+) and (π^-, K^+) in nuclear targets. The potentials in SNM for the other hyperonic species are unfortunately less precisely known. We utilized the value $U_{\Xi} = -24$ MeV reported in Ref. [291], which was calculated taking into account the few two-body Ξ capture events to two single Λ -hypernuclei observed in The High Energy Accelerator Research Organization (KEK) and The Japan Proton Accelerator Research Complex (J-PARC) [292, 293, 294]. For Σ hyperons, the value of $U_{\Sigma} = 30$ MeV was adopted. The value was obtained in Refs. [295, 296], taking into account the Σ^- spectrum measurements at KEK [297, 298]. Additionally, the Λ - Λ bond energy, $\Delta B_{\Lambda\Lambda}(\rho_0/5) = -0.67$ MeV, reported in [299, 300], which was extracted from the binding energy of $^6_{\Lambda\Lambda}\text{He}$, is used to constrain the $\sigma^* - \Lambda$ coupling, $g_{\sigma^*\Lambda}$.

However, there is a significant uncertainty associated with these values. As those uncertainties propagate at higher densities, they can influence the composition and the EoS of hypernuclear matter significantly. Therefore, we utilize two extreme models,

FSU2H*U and FSU2H*L, whose coupling constants are obtained by reproducing the hyperonic potentials to the upper and lower limits of their experimental uncertainty bands, respectively. The Λ hyperon band is pretty narrow, $U_\Lambda = [-30, -25]$ MeV, while the other hyperonic potentials have wider uncertainties: $U_\Xi = [-24, -10]$ MeV, $U_\Sigma = [10, 50]$ MeV and $\Delta B_{\Lambda\Lambda}(\rho_0/5) = [-6, 0]$ MeV [291, 296, 301, 302, 303]. Table 2.3 presents the parameter values as ratios relative to their nucleonic counterparts. For non-hidden mesons, these ratios ($R_{iY} = g_{iY}/g_{iN}$) compare their couplings to those of nucleons for the same meson. However, because of the OZI rule, σ^* and ϕ do not couple to nucleons [304]. Consequently, their coupling ratios are defined differently: $R_{\sigma^*Y} = g_{\sigma^*Y}/g_{\sigma N}$ and $R_{\phi Y} = g_{\phi Y}/g_{\omega N}$.

	ρ_0 [fm ⁻³]	E_0 [MeV]	K_0 [MeV]	$E_{\text{sym},0}$ [MeV]	L_0 [MeV]	P_{PNM} [MeV fm ⁻³]
FSU2H*	0.1505	-16.28	238.0	30.5	44.5	2.30
Constraints	0.148 ± 0.004 [280]	-16.24 [283]	240 ± 20 [305, 306, 307, 308]	31.6 ± 2.7 [284]	54 ± 8 [284]	2.45 ± 0.15 [287]

Table 2.1: Nuclear matter properties predicted by the FSU2H* model. The nuclear constraints obtained in various experimental and theoretical works are also listed.

m_σ [MeV]	m_ω [MeV]	m_ρ [MeV]	m_{σ^*} [MeV]	m_ϕ [MeV]	$g_{\sigma N}^2$	$g_{\omega N}^2$	$g_{\rho N}^2$	κ [MeV]	λ	ζ	Λ_ω
497.479	782.5	763	980	1020	102.72	169.53	197.27	4.0014	-0.0133	0.008	0.045

Table 2.2: Parameters of the model FSU2H* for the nucleon coupling constants and meson masses. The mass of the nucleon is equal to $m_N = 939$ MeV. The table was published in [259].

Y	$R_{\sigma Y}$	$R_{\sigma Y}(U)$	$R_{\sigma Y}(L)$	$R_{\sigma^* Y}$	$R_{\sigma^* Y}(U)$	$R_{\sigma^* Y}(L)$	$R_{\omega Y}$	$R_{\rho Y}$	$R_{\phi Y}$
Λ	0.6113	0.6048	0.6178	0.2812	0.2309	0.4954	2/3	0	$-\sqrt{2}/3$
Σ	0.4673	0.4085	0.5132	0.2812	0.2309	0.4954	2/3	1	$-\sqrt{2}/3$
Ξ	0.3305	0.2938	0.3325	0.5624	0.4618	0.9908	1/3	1	$-2\sqrt{2}/3$

Table 2.3: Hyperon coupling ratios to mesons (σ , ω , ρ , σ^* and ϕ) are presented, relative to the nucleon couplings. σ and σ^* couplings are model-dependent (FSU2H*, FSHU2*U (U), FSHU2*L (L)), while other ratios are consistent across models as they are related with symmetry relations to the same nucleonic core of the model. The table was published in [309].

2 Composition and EoS of hypernuclear matter

We begin this section by studying the composition of hypernuclear matter for a wide range of values of T , Y_Q , and ρ_B , covering all conditions that could be present during different astrophysical scenarios. We do the analysis utilizing the models that we have developed in the previous section - FSU2H*L and FSU2H*U. We analyze matter at three distinct temperatures (low: $T = 1$ MeV, moderate: $T = 20$ MeV, and high: $T = 80$ MeV) and three charge fractions ($Y_Q = 0.01$, $Y_Q = 0.25$ and $Y_Q = 0.5$).

The abundance of various baryons as a function of the baryon density is shown in Fig. 2.1. Solid lines correspond to calculations with the FSU2H*U model, while dashed lines correspond to those with FSU2H*L. Each row corresponds to a different temperature: $T = 1$ MeV (up), $T = 20$ MeV (middle), and $T = 80$ MeV (down), while the columns represent the different constant charge fractions: $Y_Q = 0.01$ (left), $Y_Q = 0.25$ (middle), and $Y_Q = 0.5$ (right). Before we start the analysis, it should be noted that the precise combination of the different hyperonic potentials can have an impact on species abundances. Consequently, the uncertainty density region for each hyperonic particle is only partially represented by our two extreme models, FSU2H*U and FSU2H*L. However, the uncertainty in the total amount of hyperons present in matter is captured with our two parametrizations (see Appendix B for more details).

Note that in the absence of hyperons, the composition is straightforward, dictated solely by the charge fraction. The neutron and proton partial densities are defined as $\rho_n = (1 - Y_Q)\rho_B$ and, $\rho_p = Y_Q\rho_B$, respectively, leading to fractions $x_n = (1 - Y_Q)$ and $x_p = Y_Q$. These simple relationships break down with the appearance of hyperons, where, due to the weak interaction equilibrium, the particle abundances become dependent on both temperature and baryon density.

At low temperatures, the onset density for each hyperon is highly sensitive to the hyperonic potentials. This is easily seen by inspecting the composition patterns of the Σ^- hyperons, which exhibit the largest potential uncertainties. For example, at $Y_Q = 0.01$, the significant chemical potential difference between neutrons and protons, together with the soft repulsive potential, allows Σ^- hyperons to appear even before Λ hyperons in the FSU2H*L model. Conversely, the FSU2H*U model's strongly repulsive potential delays the onset of Σ^- hyperons, which within this model appear after Λ and Ξ^- hyperons. Furthermore, at low charge fractions, the rising abundance of negative hyperons causes a corresponding increase in the proton fraction. However, this does not hold for moderate charge fractions. At low temperature and charge fractions $Y_Q = 0.25$ and $Y_Q = 0.5$, the proton abundance remains relatively stable and high, the difference between the chemical potential of the neutron and the proton is not substantial, and the emergence of negative hyperons is suppressed. Nevertheless, neutrons can always be substituted by non-degenerate, low-energy Λ hyperons, regardless of the value of Y_Q , leading to a

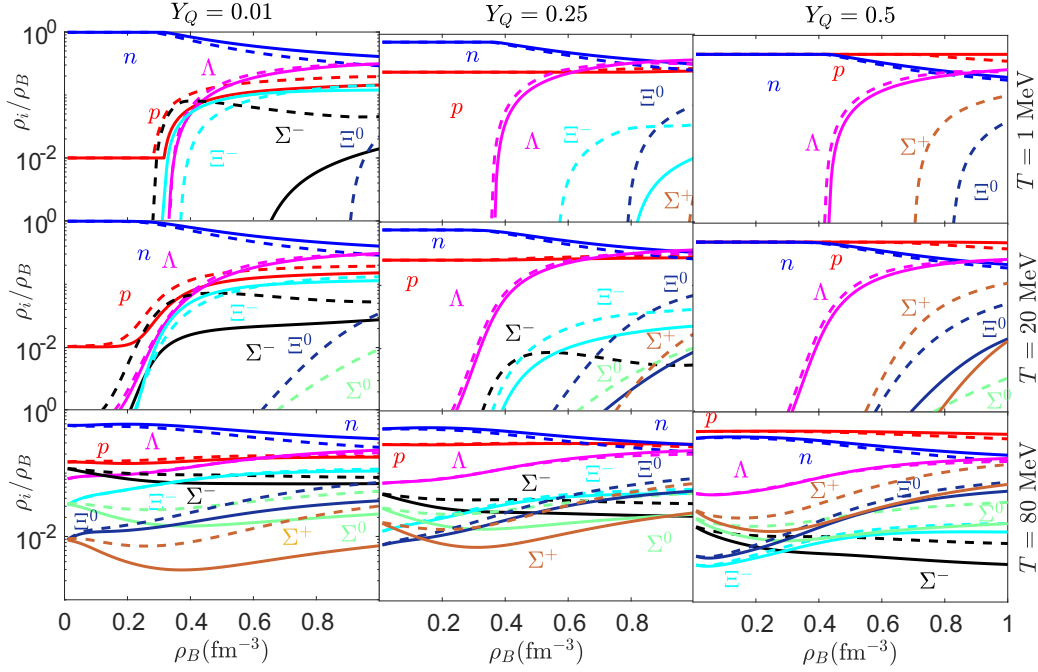


Figure 2.1: Composition of baryonic matter at constant charge fraction and constant temperature. The results obtained with the FSU2H*U model are labeled with solid lines, while the dashed lines correspond to the results obtained with the FSU2H*L model. Different species of matter are represented by each color. The figure is published in [289].

significant decrease in their abundance. The onset density at which this reduction starts is, however, charge fraction dependent.

The qualitative conclusions that we have derived at low temperatures hold on at $T = 20$ MeV, although some key distinctions emerge. First, the onset density of any hyperonic species, at any charge fraction, due to the thermal excitations, shifts toward lower densities. The onset stops being abrupt, and hyperons appear steadily in matter. Furthermore, at a given density, a larger number of hyperonic species can coexist. This is particularly true when the potentials that hyperons feel are softer (see, for example, the FSU2H*L calculations in the middle panel).

When the temperature is high enough ($T = 80$ MeV), the conclusions we have drawn for the situations at low and moderate temperatures do not hold. As it can be easily seen, all hyperons are present in matter in significant quantities, at any charge fraction, even at ultra-low densities. The uncertainty in the Σ hyperon potential that was influencing the abundance of the Σ^- hyperon at low temperature now particularly influences the abundance of the Σ^+ . As it is the only hyperon with a positive charge, its appearance plays a crucial role in reducing the proton degeneracy. The abundance predicted by the

FSU2H*L model of Σ^+ and n are nearly similar at ultra-large densities, as can be seen in the down-right panel of the figure ($T = 80$ MeV and $Y_Q = 0.5$).

The differences that exist in the composition of the matter between the different cases also affect the EoS. We will focus on inspecting three different quantities: the pressure P , the energy per particle E/A , and the entropy S/A . The results are shown in Fig. 2.2, where the legend indicates the temperature of each curve expressed in MeV. Each column represents a different charge fraction, while each row represents a different thermodynamic quantity of interest. The results at $T = 1$ MeV are labeled with blue lines, the ones at $T = 20$ MeV are marked with green lines, while the red lines represent the calculations at $T = 80$ MeV. As in the previous figure, the solid line corresponds to calculations with the FSU2H*U model, while the dashed lines label the calculations with the FSU2H*L model.

It is evident that hyperonic uncertainties have the biggest impact on the pressure. This effect may potentially outweigh the thermal contribution effect at large densities. We first look at the $P(\rho_B)$ dependence at $Y_Q = 0.01$ to illustrate this. The $T = 80$ MeV FSU2H*L curve (dashed red line) is below the $T = 1$ MeV FSU2H*U curve (solid blue line) for densities $\rho_B \gtrsim 0.6 \text{ fm}^{-3}$. The conclusion becomes more apparent as we shift towards more isospin-symmetric matter. As matter becomes proton-richer, hyperon abundance diminishes, according to Fig. 2.1. Consequently, the high-density matter becomes more degenerate, resulting in $P(\rho_B)$ curves for different temperatures that exhibit minimal variation in the stiffest model, FSU2H*U. The softer model, FSU2H*L, permits a substantial abundance of highly degenerate protons to be converted into Σ^+ hyperons. As a consequence of having more species in matter, the FSU2H*L curves, even in symmetric matter, exhibit a variation with temperature.

Regarding the density dependence of the energy per particle, Fig. 2.2 clearly demonstrates that the discrepancies between the two models are more significant at high densities and high temperatures. Furthermore, the impact is increased when the total hyperon abundance of matter is higher, as is evident from the left-middle panel ($Y_Q = 0.01$).

The entropy per particle is illustrated in the lower panels of Fig. 2.2. It is noticeable that S/A exhibits no clear, significant dependence on the hyperonic potentials. This suggests that the temperature profiles during the early evolution of stars are not much impacted by the uncertainties in the YY and YN interactions. This may seem counterintuitive, as the appearance of hyperons in matter has a significant effect on the temperature in PNSs. As detailed in [289, 310, 311], the temperature profiles computed for stars at constant S/A do, in fact, suffer substantial modifications with the appearance of hyperons, developing a plateau-like characteristic. However, the two statements are not contradictory. When energetically favorable, the most degenerate nucleons are converted into hyperons, resulting in the filling of a new Fermi sea. The reduction of the degeneracy strongly influences the entropy, leading to significantly larger values and the development

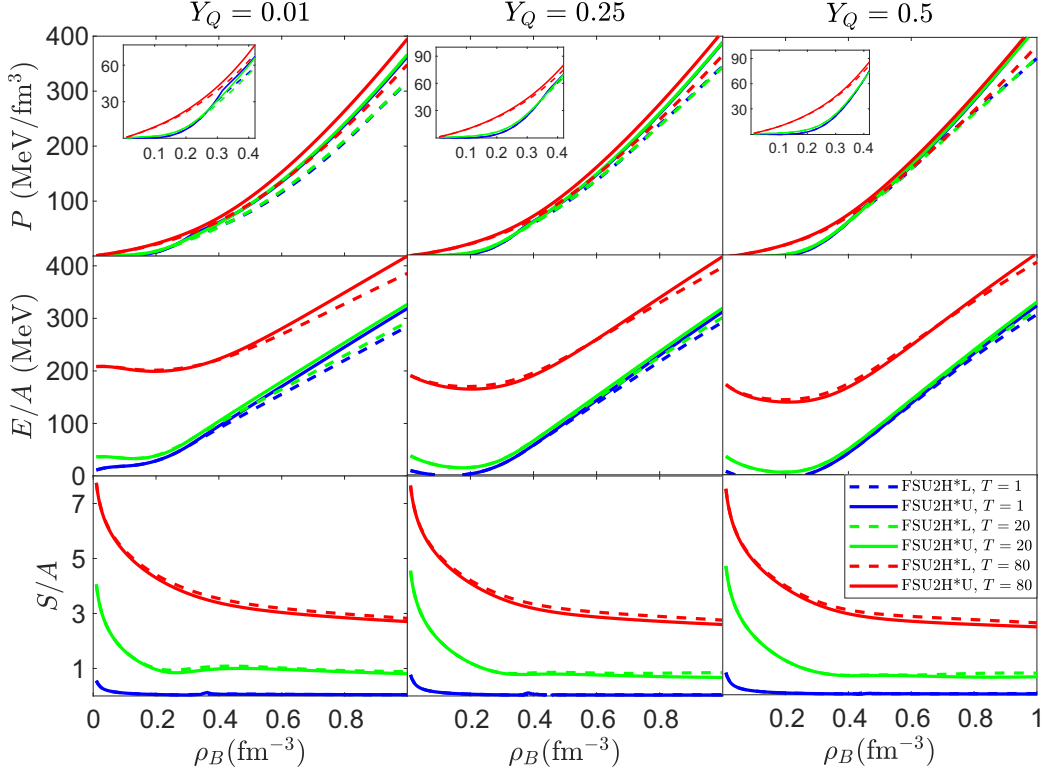


Figure 2.2: Pressure P , energy per particle E/A and entropy S/A , as functions of baryon density for hypernuclear matter at different charged fractions and temperatures. Dashed lines indicate calculations using the FSU2H*L model, and solid lines indicate calculations using the FSU2H*U model. The figure is published in [289].

of the plateau-like structure (see, for example, the temperature profiles at $S/A = 1$ in Fig. 3.1 - Chapter 3). Nevertheless, the effect of hyperonic potential uncertainty on S/A is negligible once hyperons are present.

3 Thermal quantities in β -stable NS matter

As we have previously discussed, the most difficult task of determining the EoS for NS matter at finite temperature is finding the strongly correlated baryonic EoS. In many cases the leptons, specifically electrons and muons, are generally not in full chemical equilibrium in NS matter. Consequently, their contributions are typically not directly tabulated within the EoS tables. Instead, the explicit calculation of lepton reaction rates and the solution of their respective transport equations become necessary to fully describe the thermodynamic state and evolution of the system.

However, in the highest densities in Core Collapse Supernova (CCSN) and in the subsequent evolution of PNSs, the β -equilibrium condition can be established [244], which greatly simplifies the problem. Additionally, the NS matter must always maintain charge neutrality in order to prevent destabilizing strong electric fields. This results in the following relationship between the charged lepton densities ρ_l and the charge baryon density $Y_Q \rho_B$:

$$Y_Q \rho_B = \sum_{l=e,\mu} \rho_l. \quad (2.26)$$

These two assumptions allow a NS EoS of β -stable matter to be built. Depending on the specific conditions, the neutrinos can be either trapped in the matter or they can become untrapped. In the early phases of the PNS, they are indeed trapped; their chemical potential is non-vanishing, and it is related to the chemical potential of other species through relations like:

$$\mu_n + \mu_{\nu_l} = \mu_p + \mu_l, \quad (2.27)$$

where $l = e, \mu^-$. In those phases, the total lepton numbers:

$$Y_l = \frac{\rho_l + \rho_{\nu_l}}{\rho_B} \quad (2.28)$$

are conserved. We will not consider the possibility of oscillation between the different neutrino species. In the latter stages of a PNS evolution, the mean free path of the neutrinos becomes larger than the dimensions of the object, allowing them to stream freely from the star [312]. Consequently, their chemical potential becomes zero, and the lepton numbers are no longer conserved. The chemical potentials of the muons and electrons become equal, so $\mu_e = \mu_\mu$, and they are related to the chemical potentials of the baryons through the relation:

$$\mu_n = \mu_p + \mu_{e/\mu}. \quad (2.29)$$

In all conditions of interest, leptons can be treated as non-interacting particles. Their abundance in matter is given by:

$$\rho_{l(\nu_l)} = \frac{\gamma_{l(\nu_l)}}{2\pi^2} \int_0^\infty (f_{l(\nu_l)}(k, T) - f_{\bar{l}(\bar{\nu}_l)}(k, T)) dk k^2, \quad (2.30)$$

where $f_{l(\nu_l)}(k, T)$ and $f_{\bar{l}(\bar{\nu}_l)}(k, T)$ are the Fermi-Dirac distributions for leptons and their corresponding antiparticles, respectively:

$$f_{l/\nu_l(\bar{l}/\bar{\nu}_l)}(k, T) = \left[1 + \exp \left(\frac{\sqrt{k^2 + m_{l/\nu_l(\bar{l}/\bar{\nu}_l)}^2} - \mu_{l/\nu_l(\bar{l}/\bar{\nu}_l)}}{T} \right) \right]^{-1}. \quad (2.31)$$

Here, $\mathcal{Y}_{l(v_l)}$ is the degeneracy factor. Similar to the case for baryons, it holds that $\mu_{l/v_l} = -\mu_{\bar{l}/\bar{v}_l}$. The total lepton contribution to the energy of the system [278] can be written as:

$$\varepsilon_L = \sum_{l/(v_l)} \mathcal{Y}_{l(v_l)} \int_0^\infty \sqrt{k^2 + m_{l(v_l)}^2} (f_{l(v_l)}(k, T) + f_{\bar{l}(\bar{v}_l)}(k, T)) k^2 dk. \quad (2.32)$$

Similarly, the pressure can be expressed as:

$$P_L = \sum_{l/(v_l)} \mathcal{Y}_{l(v_l)} \int_0^\infty \frac{k^2}{\sqrt{k^2 + m_{l(v_l)}^2}} (f_{l(v_l)}(k, T) + f_{\bar{l}(\bar{v}_l)}(k, T)) k^2 dk. \quad (2.33)$$

The total energy density and total pressure of the system are the sum of the lepton and baryon contributions:

$$\begin{aligned} \varepsilon &= \varepsilon_B + \varepsilon_L; \\ P &= P_B + P_L. \end{aligned} \quad (2.34)$$

The lepton entropy density (s_L) and the lepton free energy density (f_L) are derived from the universal thermodynamic relations:

$$s_L = \frac{1}{T} \left(\varepsilon_L + P_L - \sum_{l/(v_l)} \mu_{l/(v_l)} \rho_{l/(v_l)} \right); \quad (2.35)$$

$$f_L = \varepsilon_L - T s_L. \quad (2.36)$$

Finally, the total entropy and free energy of the system are also the sums of the respective baryon and lepton contributions:

$$\begin{aligned} s &= s_B + s_L; \\ f &= f_B + f_L. \end{aligned} \quad (2.37)$$

We are interested to see the thermal behavior of hypernuclear NS matter in β -equilibrium and see whether that differs with respect to the one of the usual nuclear NS matter only. In this part we will use only our nominal FSU2H* model, as the conclusions that we will draw are not sensitive to the potential that hyperons feel in the matter. However, we will also consider a nucleonic-only version of the FSU2H* by artificially prohibiting the appearance of hyperons in the matter. That can be done by setting the masses of the hyperons to be infinitely large. The comparison between those two models will tell us the difference that arises from the appearance of the hyperons. For simplicity, we will focus only on the cases where neutrinos are not present in the matter. However, as we

explicitly showed in [289], the inclusion of the neutrinos does not change the qualitative conclusions.

The main quantity that we are interested to inspect is the so-called thermal index Γ . The thermal index in β -stable matter is a function of the baryon density and the temperature and is related to the ratio of the thermal pressure with respect to the thermal energy density:

$$\Gamma(\rho_B, T) \equiv 1 + \frac{P_{\text{th}}}{\varepsilon_{\text{th}}} . \quad (2.38)$$

The thermal pressure and thermal energy density represent the difference between the pressure and the energy density at some finite temperature T and their values at zero temperature, respectively:

$$\begin{aligned} P_{\text{th}}(\rho_B, T) &= P(\rho_B, T) - P(\rho_B, 0); \\ \varepsilon_{\text{th}}(\rho_B, T) &= \varepsilon(\rho_B, T) - \varepsilon(\rho_B, 0) . \end{aligned} \quad (2.39)$$

The thermal index in NS matter is conceptually analogous to the familiar adiabatic index found in the ideal gas law. For an ideal gas, the adiabatic index is a constant, equal to 5/3 for non-relativistic fermions and 4/3 for ultra-relativistic fermions, and is independent of temperature and density.

Motivated by the ideal-gas behavior, the constant thermal index approximation in merger simulations has been extensively used. To simplify calculations while at the same time capturing essential physics, the community has combined well-established cold matter EoSs with an additional ad hoc thermal correction based on the thermal index [313, 314, 315]. Within this approach, the thermal energy density is treated as an independent variable. The thermal index, is replaced by a constant Γ , typically taking a value between 1.5 and 2.0. The method yields a finite temperature EoS, $P(\varepsilon)$ if the zero-temperature pressure P_0 and energy density ε_0 are known:

$$P(\varepsilon) = (\Gamma - 1)(\varepsilon - \varepsilon_0) + P_0. \quad (2.40)$$

However, NS matter is fundamentally different than the ideal-gas systems, comprising of strongly interacting baryons and non-interacting leptons. The strong interaction leads to deviations from ideal gas behavior, making the thermal index a quantity that is not a simple constant. While computationally efficient, the simplified approach can introduce inaccuracies under the extreme conditions of BNS mergers [316, 317].

We start our analysis by inspecting the behavior of P_{th} and ε_{th} as functions of baryon density within our FSU2H* model (Fig. 2.3). Calculations for three different temperatures are shown: $T = 10$ MeV as black curves, $T = 30$ MeV as green curves, and $T = 50$ MeV as red curves. The calculations where hyperons are considered in matter are labeled with

dashed lines, while the ones where they are artificially prevented from appearance are labeled with solid lines.

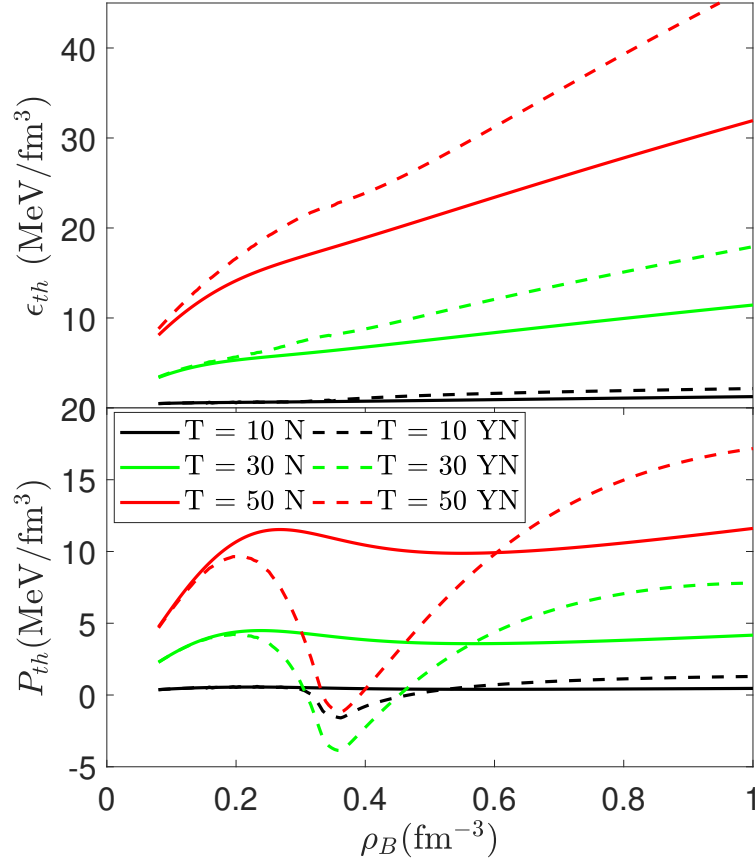


Figure 2.3: Thermal energy density and thermal pressure as functions of the baryonic density for different temperatures, expressed in MeV units. The solid lines label the cases where hyperons are not present in the matter, while the dashed ones represent the calculations where hyperons are taken into account. The figure was published in [289].

We can see that the appearance of hyperons increases the value of ϵ_{th} . The increase is more important at higher temperatures. However, the effect that hyperons have on P_{th} is significantly greater. Near the hyperon appearance threshold at $T = 0$ (around 0.35 fm^{-3}), P_{th} experiences a noticeable dip, deepening with increasing temperature. Notice that the lowering of the thermal pressure starts at earlier densities when the temperature is higher, suggesting that this effect is directly related to the amount of thermal hyperons that are produced in the matter. Indeed, we see that the dips at higher temperatures are wider as

well. At higher densities, P_{th} recovers and surpasses the values obtained in hyperon-free matter. At those densities, the difference between the hyperon abundance at zero and at finite temperature is reduced. We will come back to this point in Chapter 4.

This intricate behavior of P_{th} directly impacts the thermal index, as shown in Fig. 2.4 for various temperatures, labeled with different colors. The hyperonic EoS (dashed lines) exhibits a sharp dip, corresponding to the dip that P_{th} has. A greater increase in ϵ_{th} causes this dip to become less noticeable as temperature rises, and it becomes wider when hyperons emerge at lower densities. It is interesting to see that at higher densities, the thermal indices of hypernuclear matter are higher than those of solely nucleonic matter. This discrepancy results from the abundance of nucleons in nucleonic matter moving toward the degenerate ultra-relativistic regime ($\Gamma = 4/3$). In those conditions, the scalar fields increase dramatically, which results in a significantly lower effective mass (see Eq. (2.14)). The behavior is relatively temperature-independent, typical of degenerate systems. However, in hypernuclear matter, the presence of hyperons causes the higher momentum nucleonic states to be emptied and lower hyperonic momentum states to be filled, both of which help to shift the thermal index toward higher values.

To better quantify the potential errors introduced by using a constant thermal index when calculating the thermal effects, in Fig. 2.5 we display the thermal pressure against the thermal energy density. The calculations with the nucleonic-only version of the model are marked with solid lines, while the calculations with the full FSU2H* model are labeled with dashed lines. Two different temperatures are considered: $T = 30$ MeV - green lines and $T = 50$ MeV - red lines. While only two temperatures are shown, a broader and denser temperature grid would produce families of nucleonic and hyperonic curves spanning the different regions of the $P_{\text{th}} - \epsilon_{\text{th}}$ plane. Two common ($\Gamma = 1.5$ and $\Gamma = 2$), simplistic linear approximations employed in merger simulations are shown by the dotted black lines, which draw attention to the differences from our consistent calculations. Indeed, Fig. 2.5 reveals that the relationship between P_{th} and ϵ_{th} is far more intricate than the linear approximation commonly employed, even in the case where nucleons are the only baryons considered in matter. In fact, for a specific value of ϵ_{th} , P_{th} can vary significantly with a change of temperature, contradicting the temperature-independent thermal correction implied by assuming a constant thermal index. For example, at $\epsilon_{\text{th}} = 10$ MeV/fm³, the nucleonic thermal pressure at the two temperatures differs by an amount of 3 MeV/fm³. The difference is larger than 6 MeV/fm³ in the case where hyperons are also present in the matter. While a constant index of $\Gamma = 1.5$ seems to approximate the nucleonic behavior on average, the presence of hyperons turns the use of any constant value invalid. The hyperonic calculations span a much wider region in the $P_{\text{th}} - \epsilon_{\text{th}}$ plane, highlighting the inadequacy of a simplified, constant thermal index approach.

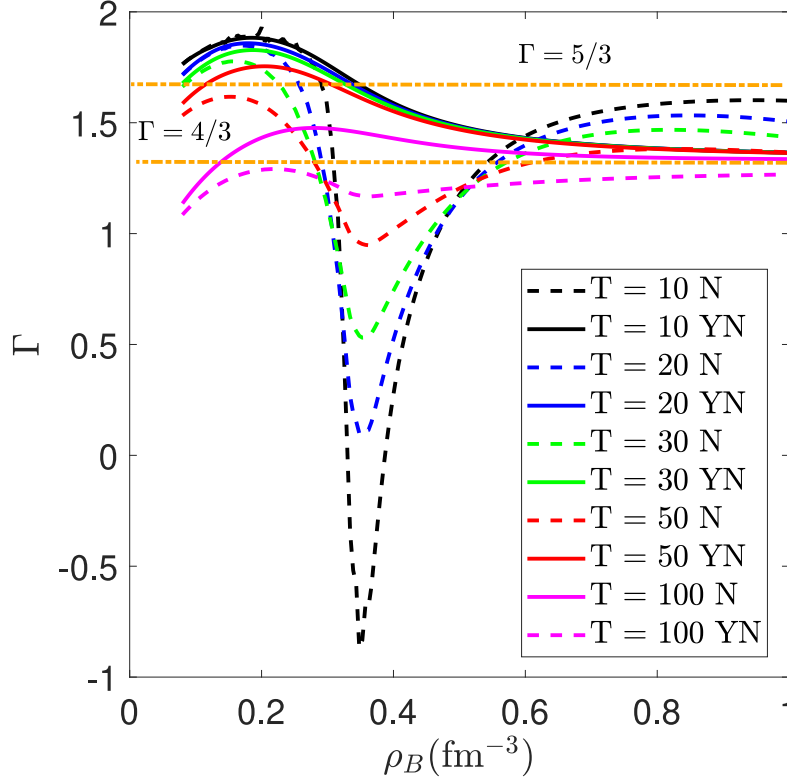


Figure 2.4: Thermal index as a function of the baryon density for various temperatures, expressed in MeV units. The calculations performed without hyperons are marked with solid lines, while the ones that do account for the appearance of hyperons are labeled with dashed lines. The horizontal dashed-dotted orange lines correspond to the thermal index of a non-relativistic ($\Gamma = 5/3$) and ultra-relativistic ($\Gamma = 4/3$) diluted ideal gas, respectively. The figure was published in [289].

Acknowledging the inaccuracy of assuming a constant thermal index for approximating thermal effects, while also recognizing the computational costs of complex relativistic simulations, we present parameterizations of Γ for both nucleonic and hyperonic neutrinoless β -stable matter. These parameterizations are valid for temperatures ranging from 10 MeV to 100 MeV and baryon densities from 0.08 fm^{-3} to 1.0 fm^{-3} . For matter composed solely of nucleons as the baryonic degree of freedom, the behavior of the thermal index is simpler and can be accurately described by the following form:

$$\Gamma(\rho_B, T) = a_1 + b_1 e^{-c(T)\rho_B^2} + d_1 e^{-e(T)\rho_B}, \quad (2.41)$$

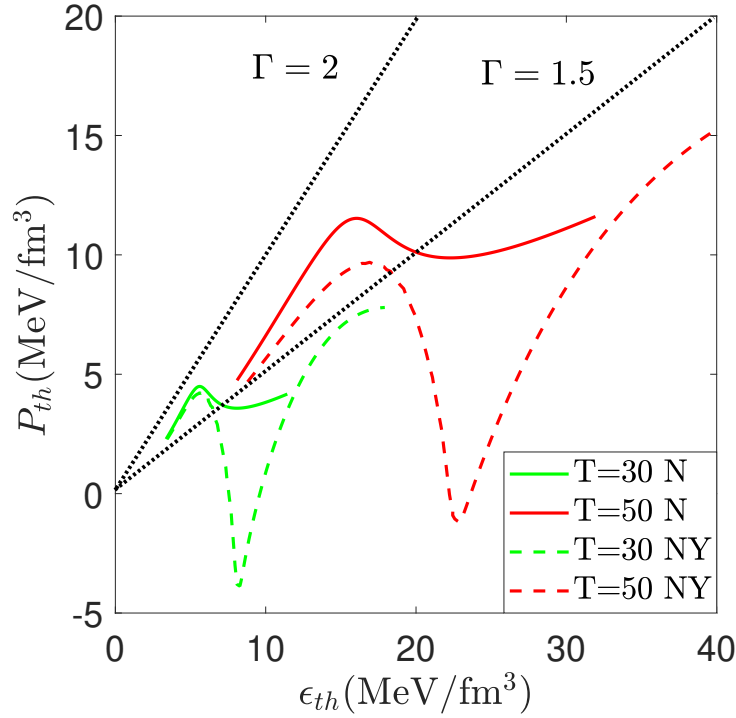


Figure 2.5: P_{th} as a function of ϵ_{th} for various temperatures, expressed in MeV units. The calculations without hyperons are marked with solid lines, while the ones with hyperons are labeled with dashed lines. The dotted lines represent two linear approximations of the $P_{th}(\epsilon_{th})$ relation ($\Gamma = 2$ -upper line and $\Gamma = 1.5$ -lower line). The figure was published in [289].

where $c(T)$ and $e(T)$ are linear functions of the temperature:

$$\begin{aligned} c(T) &= c_1 T + c_2; \\ e(T) &= e_1 T + e_2. \end{aligned} \quad (2.42)$$

However, when hyperons are present in matter, the behavior of Γ is much more complex, and simple expressions like Eq. (2.41) are not sufficient to reproduce the results. Instead, we propose the following functional form:

$$\Gamma(\rho_B, T) = \frac{a(T)\rho_B + b(T)}{(\rho_B - \rho_{Bc})^2 + c(T)\rho_B + d(T)} + e(T)\rho_B + f(T)\sqrt{\rho_B} + g(T), \quad (2.43)$$

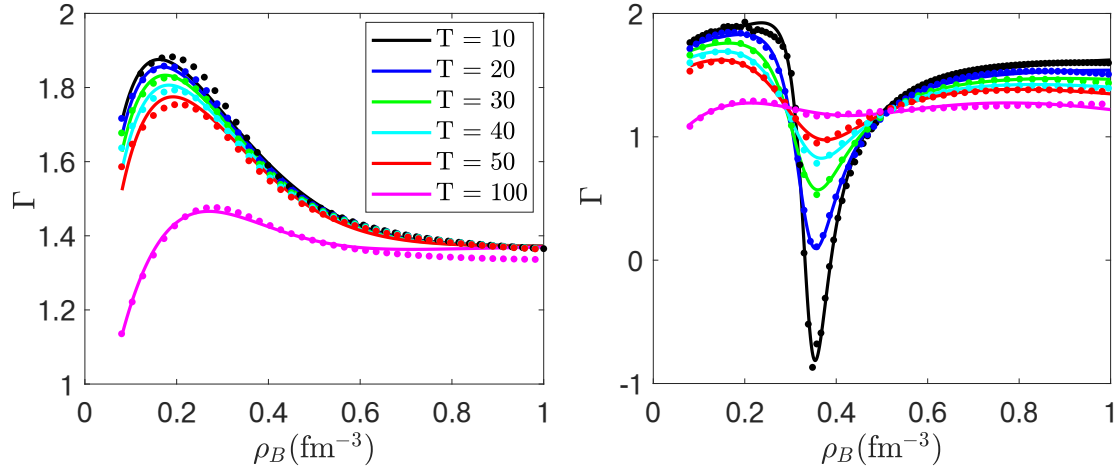


Figure 2.6: $\Gamma = \Gamma(\rho_B, T)$ fit as a function of the baryonic density ρ_B for the neutrino-free case, without (left panel) and with hyperons (right panel). Dots are representing the exact calculations, while the solid lines are the results from the fits. The temperatures shown are in MeV units. The figure was published in [289]

where $\rho_{B_c} = 0.3715 \text{ fm}^{-3}$, and the temperature-dependent functions depend either linearly or quadratically on the temperature:

$$\begin{aligned}
 a(T) &= a_1 T + a_2; \\
 b(T) &= b_1 T + b_2; \\
 c(T) &= c_1 T^2 + c_2 T + c_3; \\
 d(T) &= d_1 T^2 + d_2 T + d_3; \\
 e(T) &= e_1 T^2 + e_2 T + e_3; \\
 f(T) &= f_1 T + f_2; \\
 g(T) &= g_1 T^2 + g_2 T + g_3.
 \end{aligned} \tag{2.44}$$

In Fig. 2.6 we demonstrate that this approach accurately reproduces the value of Γ regardless of the presence of hyperons in the matter. Indeed, the solid curves, which represent the values of the parametrizations, closely follow the corresponding dots, which are the exact calculations using our FSU2H* model.

For practical use, the values of the parameters of the fits are given in Table 2.4 and Table 2.5. The units of each coefficient are such that Γ remains a dimensionless quantity.

i	a_i	b_i	c_i	d_i	e_i
1	1.37	$6.89 \cdot 10^{-1}$	$1.69 \cdot 10^{-3}$	1.47	$-1.41 \cdot 10^{-1}$
2	/	/	6.59	/	20.29

Table 2.4: Thermal index fit coefficients when nucleons are the only baryon present in the β -stable matter. The units of each coefficient are such that Γ remains a dimensionless quantity. The table was published in [289].

i	a_i	b_i	c_i	d_i	e_i	f_i	g_i
1	$-2.02 \cdot 10^{-4}$	$1.98 \cdot 10^{-2}$	$-1.71 \cdot 10^{-5}$	$1.24 \cdot 10^{-5}$	$1.11 \cdot 10^{-4}$	$3.21 \cdot 10^{-2}$	$-6.64 \cdot 10^{-5}$
2	$-6.14 \cdot 10^{-2}$	$-1.26 \cdot 10^{-4}$	$1.83 \cdot 10^{-3}$	$-6.22 \cdot 10^{-4}$	$-3.63 \cdot 10^{-2}$	$2.60 \cdot 10^{-1}$	$-5.13 \cdot 10^{-3}$
3	/	/	$1.56 \cdot 10^{-2}$	$-4.83 \cdot 10^{-3}$	$7.15 \cdot 10^{-2}$	/	1.48

Table 2.5: Thermal index fit coefficients when hyperons are also present in the β -stable matter. The units of each coefficient are such that Γ remains a dimensionless quantity. The table was published in [289].

Chapter 3

Hyperonic effects on neutron star observables

In the previous chapter, we explored how hyperons can be incorporated as a degree of freedom in a RMF theoretical framework to calculate the EoS of dense matter at both zero and finite temperature. We observed that their presence significantly alters the properties of dense matter, particularly the pressure. The uncertainties in determining the model parameters were also discussed.

This chapter investigates how these effects translate into the predictions derived from our model for the global properties of NSs. We will focus on properties that can be observationally determined, such as the mass and the radius of NSs, their tidal deformability, and the moment of inertia. Since the effects on these observables depend on the amount of hyperons present in the stars, we also perform calculations of the strangeness number. Whenever possible, we will compare our predictions with the observational constraints obtained over the past decade. As these observations are coming from cold stars, we first compute the properties at zero temperature for NSs in neutrinoless beta-equilibrium. All calculations are performed with our three parametrizations: FSU2H*, FSU2H*L, and FSU2H*U in order to cover the softest and the stiffest EoS allowed by the hyperon uncertainties. In Appendix B a more thorough study, allowing for the interplay between the values of the hyperon potentials, is presented. In this way we can capture the propagation of the uncertainty of model parameters to the predictions of the global properties of the star.

Although there are currently no measurements of the global properties of NSs at finite temperature, we also calculate them for stars during their PNS evolution using our three parameterizations. For simplicity, we will assume that the β -equilibrium holds in all cases that we consider. Following the approach of Ref. [273], the conditions in these hot

stars, particularly the temperature and lepton fraction profiles, are motivated by results from sophisticated relativistic simulation. We model the PNS evolution in three distinct stages:

- *Early Evolution/Deleptonization Phase:* For the first seconds after the supernova explosion, we use temperature and lepton profiles inspired by simulations that account for neutrino transport. Specifically, our profiles are motivated by the results presented in Ref. [245] ($t = 1$ s, Pascal curve) and by results obtained with the FSU2H EoS, with the microphysics described in a series of works by [318, 319, 320, 321, 322] ($t = 1$ s, Fischer curve). The neutrinos in this phase are considered to be trapped. Both profiles are characterized by a high temperature in the low-density region and by a high lepton fraction at all densities.
- *Intermediate Stage:* Approximately 5 seconds after the explosion, we adopt the conditions from Ref. [323] where it was shown that convective motions could lead to isentropic and isoleptonic profiles much faster than previously assumed ($t = 5.1$ s, Pascal curves). The neutrinos are still trapped in matter.
- *Later Phases:* Assuming isentropic neutrino-free matter, we reproduce the conditions for the later stages of the PNS evolution, approximately tens of seconds after the explosion ($S/A = 1, 2$ and 3 curves). Matter in the high-density region is lepton poor. As pointed out in Ref. [324], these stars could be relevant for black-hole formation in a failed core collapse supernova.

The graphic representation of all profiles is given in Fig. 3.1.

We stress that our analysis is not meant to take the place of the full, intricate relativistic simulations. Rather, these findings should be interpreted as a prediction of the possible propagation of the hyperon effects on the EoS to the hot star observables. The findings we present here should hold for other phenomena with comparable conditions, such as BNS mergers, even if the conditions are modeled after those seen in PNSs.

1 Mass and radius of cold and hot NSs

The main characteristics of a NS are its mass (M) and radius (R). Due to the extreme density of NSs, their hydrostatic equilibrium cannot be theoretically described by classical Newtonian gravity. Instead, a GR treatment is required. The hydrostatic equilibrium equation for a spherically symmetric, non-rotating star, derived from Einstein's equations, is known as the Tolman-Oppenheimer-Volkoff (TOV) equation [18]

$$\frac{dP}{dr} = -\frac{\epsilon m}{r^2} \left(1 + \frac{P}{\epsilon}\right) \left(1 + \frac{4\pi P r^3}{m}\right) \left(1 - \frac{2m}{r}\right)^{-1}. \quad (3.1)$$

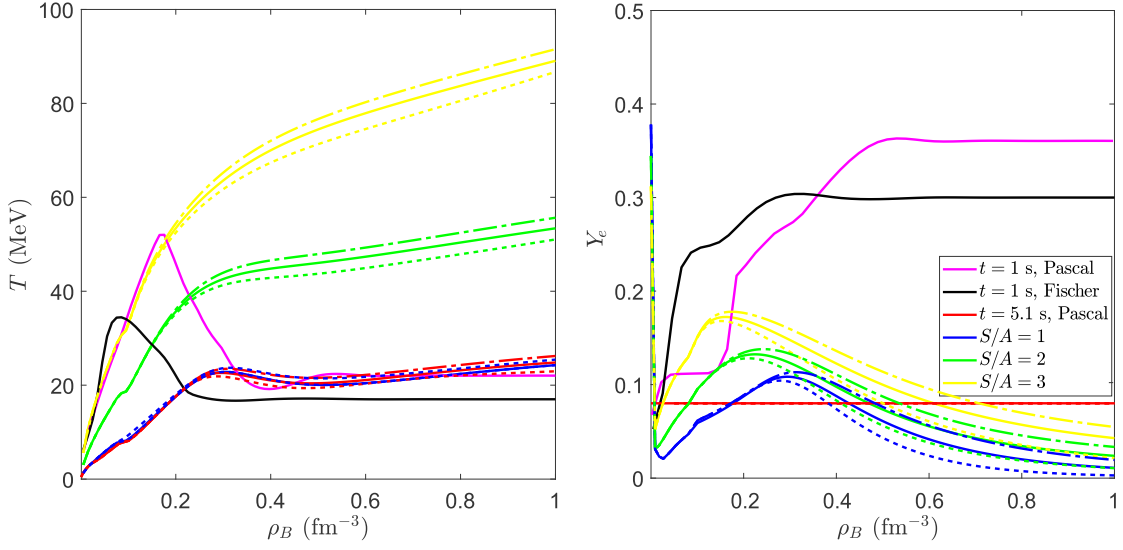


Figure 3.1: Temperature (left plot) and electron fraction profiles (right plot) of the stars used in our calculations. Solid lines represent the results with the FSU2H* model, dashed lines represent the ones obtained with the FSU2H*L model, and dashed-dotted lines are results with the FSU2H*U model. Please read the text for detailed motivation of the shown profiles. For all three models, the simulation-inspired profiles at $t = 1$ s are identical. The figure was published in [309].

Eq. (3.1) has to be supplemented for the mass m enclosed up to a radius r , which in differential form is:

$$\frac{dm}{dr} = 4\pi r^2 \varepsilon. \quad (3.2)$$

The system of equations is closed by specifying an EoS, i.e. the relation between the pressure and the energy density

$$P = P(\varepsilon). \quad (3.3)$$

In order to solve the system of equations, a central density $\varepsilon_c = \varepsilon(r=0)$ is initially fixed. The equations are then integrated outwards from the stellar center $m(r=0) = 0$, using a small radial step. The integration continues until the stellar surface is reached, which is defined as the point where the pressure becomes zero ($P(r=R) = 0$). Each value of ε_c uniquely determines the stellar mass and radius of a star for a given EoS. By varying the initial central energy density, one can generate a family of stars with different masses and radii. The curve on which those mass-radius points lie is known as the mass-radius relation $M(R)$. The maximum mass that can be produced is an important feature of each EoS. We remark that stars with central densities greater than that of the maximum-mass star are unstable to gravitational collapse and, therefore, cannot exist in nature [325].

Before presenting the $M(R)$ of cold β equilibrated NSs, obtained using the FSU2H* model and its extreme parametrizations, it is important to note that the model only describes the high-density EoS. To accurately determine the global properties of NSs, the high-density EoS must be matched with a low-density EoS. In Ref. [262], the inner crust EoS within the FSU2H model at $T = 0$ was computed, which provides a unified EoS description of the star's core-inner crust. Furthermore, we utilize the commonly used Baym-Pethick-Sutherland (BPS) EoS [326] for the outer crust, which is well-constrained by data from nuclear physics.

Fig. 3.2 presents the $M(R)$ relationship for the FSU2H* (green curve), FSU2H*L (blue curve), and FSU2H*U (red curve) parametrizations. The predictions are compared against observational constraints, including $\gtrsim 2M_\odot$ pulsar mass measurements [113, 114, 115, 116, 327] and mass-radius determinations from NICER [127, 128, 129, 130].¹ It is important to note that, similar to other hyperonic EoSs, our developed EoSs are inconsistent with the mass ($2.35 \pm 0.17 M_\odot$) of the heavy "black widow" [328, 329] pulsar PSR J0952-0607 [327]. However, the Shapiro delay, which is used to determine the mass of the other heavy NSs, is more accurate than the Keck-telescope spectrophotometry method used to assess the mass of this pulsar. Furthermore, the pulsar PSR J0952-0607 is the fastest one currently known, which can also increase the maximum mass that can be supported before the star collapses into a black hole.

We concentrate on NSs more massive than $1.6M_\odot$, where the influence of hyperons is most noticeable. The original FSU2H* model successfully satisfies all other $M(R)$ constraints. However, the FSU2H*L model struggles to reach a mass of $2M_\odot$; its soft hyperonic potentials significantly soften the EoS, reaching only $1.92M_\odot$. In contrast, the FSU2H*U model, because of the stiffer nature, achieves a higher $M_{\max}$ of $2.06M_\odot$.

While hyperonic uncertainties notably affect $M_{\max}$, their impact on the corresponding radii of stars at the maximum mass is surprisingly minimal. For example, FSU2H*U predicts a radius of approximately 12.1 km for its maximum mass star, whereas FSU2H*L predicts roughly 11.9 km. Nevertheless, for the most massive stars, hyperonic uncertainties significantly influence the radius at a given mass. The uncertainty in those cases can be as high as 10%.

Now let us focus on the $M(R)$ relation of NSs at finite temperature. Before discussing the results, several important considerations must be addressed. Similarly to the zero-temperature scenario, cluster structures may form when temperatures are below approximately 15 MeV and densities are less than half the nuclear saturation density. Note that the precise boundary of this transition depends on temperature, density, and the specific theoretical model used, as detailed in Ref. [271]. The appropriate theoretical

¹While writing the thesis, the NICER analysis of the pulsar PSR J0614-3329 was also published [133]

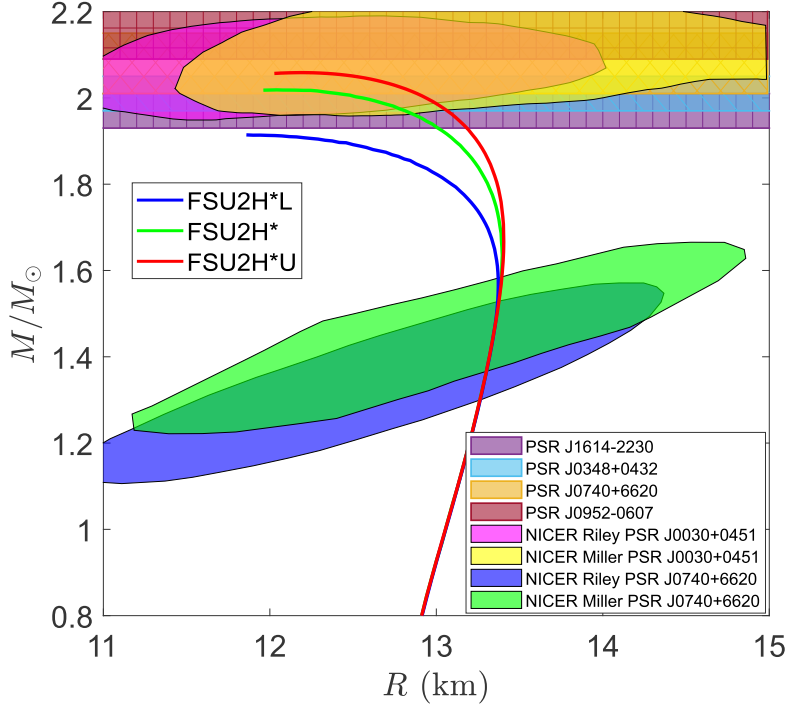


Figure 3.2: $M(R)$ curves of cold, β -equilibrated NSs, calculated using the FSU2H* family models. The observational constraints obtained from $2 M_{\odot}$ observations [113, 114, 115, 116, 327] and NICER measurements of radii within a 1σ error band [127, 128, 129, 130] are also shown. The figure was published in [309].

framework for these conditions is the extended Nuclear Statistical Equilibrium (NSE) [271, 330, 331, 332, 333, 334, 335, 336, 337, 338, 339, 340].

In this work, we use the low-density part of the EoS presented in Ref. [271], which is accessible through the CompOSE database [244, 341, 342]. The low-density EoS and our high-density EoSs are smoothly connected, and unified, so finite-temperature EoSs are obtained. While this approach may introduce some inaccuracies, particularly for the radii of the least massive stars, it provides a reasonable description for the sake of our discussion. Furthermore, determining the stellar radius at finite temperature requires a different approach than in the zero-temperature case. At $T = 0$, the radius is determined as the point where the pressure drops to zero. However, this method becomes problematic at finite temperature, due to the non-vanishing thermal pressure contribution at any density. To overcome this, we adopt the procedure outlined in Ref. [311]. We first identify the low-density region of the stars where the logarithms of pressure, energy density, and baryon density exhibit a linear relationship. This linear behavior is then used to extrapolate the EoS to even lower densities. This method provides a reasonable accuracy (within 20%) for the calculated stellar radii, particularly for low temperature

and entropy per nucleon ($S/A \leq 3$). For a more comprehensive discussion, we refer the reader to Refs. [311, 343].

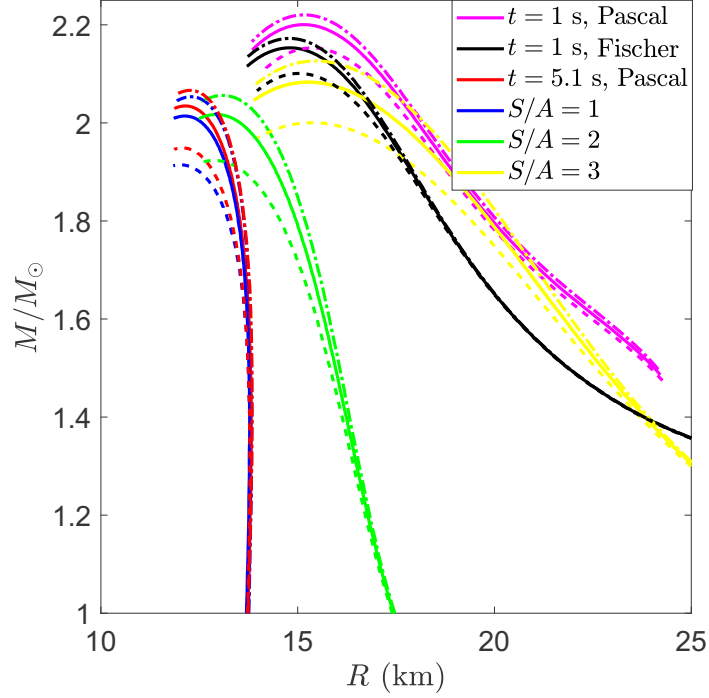


Figure 3.3: $M(R)$ relations of hot stars with temperature and density profiles as described in the text. Results from our three models are shown: FSU2H* model (solid lines), FSU2H*L (dashed lines), and FSU2H*U (dashed-dotted lines). The figure was published in [309].

Fig. 3.3 shows the $M(R)$ relation for hot stars with the temperature and lepton fraction profiles as shown in Fig. 3.1. The calculations performed with the FSU2H* models are labeled with solid lines, the ones with FSU2H*L are marked with dashed lines, while the ones with the FSU2H*U model are labeled with dash-dotted lines. Each color represents a different profile, which corresponds to the different evolution phases of the life of PNSs. From the figure, several conclusions can be drawn. The curves for the early evolutionary profiles (purple and black curves) are positioned as the rightmost in the figure, indicating a lower compactness. This is attributed to the stiffening of the EoS caused by the high temperatures present at lower core densities during these initial stages. The $M(R)$ curves of the $t = 1$ s, Pascal profile (purple curves), and the $S/A = 3$ profile (yellow curves) can be compared to better visualize this impact. Regardless of the stiffness of the selected hyperonic model or the mass of the star, their corresponding $M(R)$ curves are strikingly comparable, even though their temperature profiles only line up to the saturation density. This is further corroborated by the finding that, for densities greater than $\rho_B > 0.3 \text{ fm}^{-3}$,

the temperature profiles for $S/A = 1$ (blue curves) and $t = 1$ s, Pascal (purple curves) are fairly similar, yet their $M(R)$ curves behave very differently. This discrepancy implies that the compactness of the star is significantly influenced by the EoS at lower densities.

Another interesting feature can be deduced if one observes the profiles that describe cooler scenarios. For instance, the $t = 5.1$ s, Pascal profile (red curves), which corresponds to a constant lepton fraction case, predicts a higher maximum mass NS than the hyperon-rich profile $S/A = 1$ (blue curves). This effect is especially pronounced when hyperons experience less repulsion (dashed lines).

Similar to the zero-temperature case, we can conclude that the uncertainties in the hyperon potentials primarily impact the $M(R)$ relation for the most massive stars. In these stars, the central density is highest and the model results are specially sensitive to the particular hyperonic potential. The high temperature further amplifies these effects. Hence, the radii of the most massive stars calculated with our extreme parametrizations differ relative to one another by as much as a few tens of percent.

2 Tidal deformability of cold and hot NSs

The merger of NSs is one of the primary sources of GW signals. The existing detectors, like Advanced LIGO and Virgo [344, 345], are especially sensitive to the inspiral phase of the merging. Prior to the final merger, the binary emits a characteristic GW signal that gradually increases in frequency and amplitude, often described as a "chirp." Analysis of the waveform allows for a highly precise measurement of a parameter known as the chirp mass,

$$\mathcal{M} = \frac{(m_1 m_2)^{3/5}}{(m_1 + m_2)^{1/5}}, \quad (3.4)$$

where m_1 and m_2 are the masses of the two NSs. The intense gravitational forces at play during the late stages of the inspiral induce tidal deformations of the NSs. These deformations are quantified by the tidal deformability parameter (λ), which can also be extracted from the gravitational wave signal. Theoretically, the tidal deformability can be computed as

$$\lambda = \frac{2}{3} k_2 R^5, \quad (3.5)$$

where k_2 is the quantity often referred to as the tidal Love number. The Love number itself is a dimensionless quantity that characterizes the susceptibility of a star to be deformed by an external tidal field. By defining the compactness of the star $\beta = M/R$, k_2 can be

determined by the following robust equation [346, 347, 348, 349]

$$k_2(\beta, y_R) = \frac{8}{5}\beta^5(1-2\beta)^2[2-y_R+2\beta(y_R-1)] \\ / \left\{ 2\beta(6-3y_R+3\beta(5y_R-8))+2\beta^2[13-11y_R+\beta(3y_R-2)+2\beta^2(1+y_R)] \right. \\ \left. + 3(1-2\beta)^2[2-y_R+2\beta(y_R-1)]\log(1-2\beta) \right\}. \quad (3.6)$$

In the last equation, we used the notation $y_R = R \frac{H'(R)}{H(R)}$, where the function $H(r)$ is the solution of the following differential equation that describes the star's response to tidal forces [349]:

$$H''(r) + H'(r) \left[\frac{2}{r} + e^{\lambda(r)} \left(\frac{2m(r)}{r^2} + 4\pi r(p(r) - \varepsilon(r)) \right) \right] + H(r)Q(r) = 0, \quad (3.7)$$

and the function Q is given by

$$Q(r) = 4\pi e^{\lambda(r)} \left(5\varepsilon(r) + 9p(r) + \frac{p(r) + \varepsilon(r)}{c_s^2(r)} \right) - 6\frac{e^{\lambda(r)}}{r^2} - (v'(r))^2 \frac{e^{\lambda(r)}}{r^2}, \quad (3.8)$$

where c_s denotes the speed of sound. For spherical stars, the metric functions are given by

$$e^{\lambda(r)} = \left[1 - \frac{2m(r)}{r} \right]^{-1}, \quad v'(r) = 2e^{\lambda(r)} \frac{m(r) + 4\pi p(r)r^3}{r^2}. \quad (3.9)$$

The system of equations is closed only when the EoS is known and the proper boundary condition ($y(r=0) = 2$) is set. The assumptions with respect to the matching of the low- and high-density EoS, which we made in the previous section, will also be applied in these calculations.

In Fig. 3.4, we present the dimensionless tidal deformability ($\Lambda = \lambda/M^5$) of a cold NS as a function of its mass, plotted on a logarithmic scale. We also include existing constraints outlined in Ref. [350], extracted from the GW170817 event [351, 352]. Specifically, the tidal deformability, along with the NSs masses, spins, and other source properties, was constrained simultaneously using Bayesian parameter estimation. The method compares a waveform model to the observed data, resulting in a likelihood function across the model's parameters. The two different bands correspond to results obtained using a specific set of EoS [353] or rely on universal relations [354]. It is important to note that current constraints cover only the low- to medium-mass range of NSs, where, within our model, hyperons are not yet expected to be present.

For most of the mass range, the tidal deformability exhibits an exponentially decreasing trend. However, this trend experiences a sharp transition at a certain mass, beyond which the tidal deformability begins to lower its value much faster. A comparison with Fig. 3.2 reveals that the transition occurs at the place where a small change in the mass of the star yields a significant change in the radius (around $1.9M_{\odot}$ for both FSU2H* and FSU2H*U and around $1.7M_{\odot}$ for FSU2H*L). In other words, the tidal deformability is particularly sensitive to the change of the radius [349, 355].

We find that the tidal deformability of the most massive stars is strongly affected by the uncertainty of the hyperon interactions, especially for masses $M \gtrsim 1.8M_{\odot}$. A star with a mass of $M \approx 1.92M_{\odot}$ can be used to measure this effect. The tidal deformability for this mass is predicted by the FSU2H*U model to be $\Lambda \approx 95$, but the FSU2H*L model predicts a significantly lower value of $\Lambda \approx 34$. Therefore, we conclude that the tidal deformability of very massive, cold NSs is influenced by the hyperonic uncertainties.

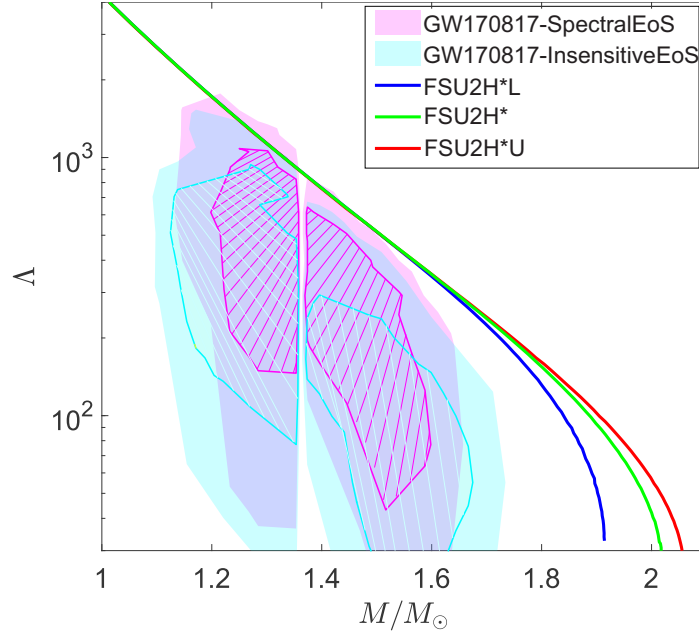


Figure 3.4: $\Lambda(M)$ relation computed with our three parametrizations together with the observational constraint on the tidal deformability using the universal relations (EoS insensitive) and the spectral EoS obtained from the GW170817 event [351, 352], presented in Ref. [350]. The hatched regions represent the values at 68% confidence level, while the filled boundaries give the values at 95% confidence level. The figure was published in [309].

Next, we analyze Fig. 3.5 that illustrates the tidal deformability of hot stars. We observe that the behavior of Λ is directly related to the temperature profile at low densities;

specifically, higher temperature values in this region yield higher Λ . This is due to the sensitivity of tidal deformability to the radius of the NSs. As shown in the previous section, a high temperature in the low-density region increases the radius of the star, which shifts the tidal deformability to larger values. Regarding the effect of hyperonic uncertainties, we find that stiffer interactions (FSU2H*U) tend to increase the tidal deformability, while softer interactions (FSU2H*L) tend to decrease it. The discrepancy between the parametrizations is more pronounced for higher-mass stars ($M > 1.7 M_\odot$), and it is especially emphasized for the most massive stars, where the prediction between the parameterizations can differ by almost an order of magnitude.

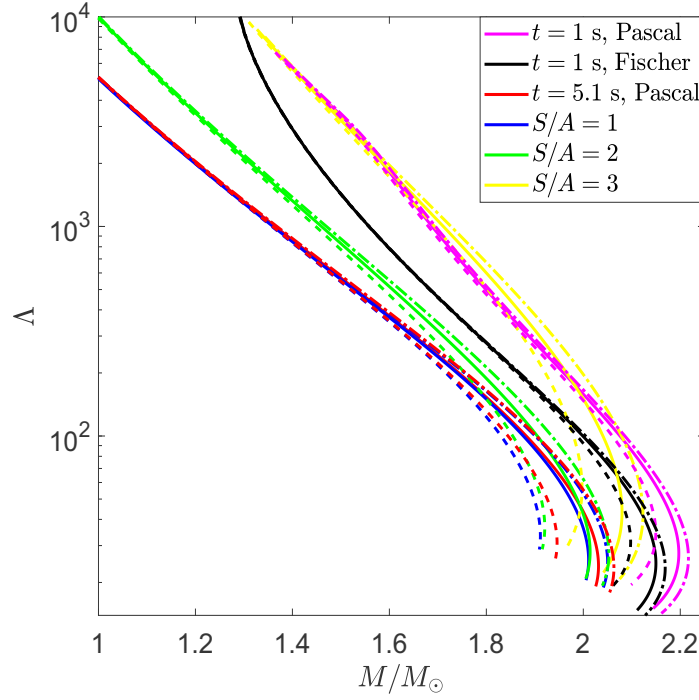


Figure 3.5: $\Lambda(M)$ relations of hot stars with temperature and lepton fraction as described in the text. Solid lines represent the results with the FSU2H* model, dashed lines represent the FSU2H*L model, and dashed-dotted lines are results with the FSU2H*U model. The figure was published in [309].

3 Moment of Inertia of cold and hot NSs

Similar to the mass-radius relation and the tidal deformability, the moment of inertia can also be computed using a GR framework. By taking the metric functions that correspond to spherical stars and assuming a slow-rotation approximation (Hartle approach), the

following equation for the moment of inertia can be derived:

$$I = \frac{w(R)R^3}{6 + 2w(R)}. \quad (3.10)$$

The radial function $w(r)$ is a solution to the following first-order differential equation:

$$\frac{dw}{dr} = 4\pi \frac{(\varepsilon + p)(4 + w)r}{1 - 2\frac{m}{r}} - \frac{w}{r}(3 + w), \quad (3.11)$$

with an initial condition $w(0) = 0$ [243, 356]. The behavior of the cold NS moment of

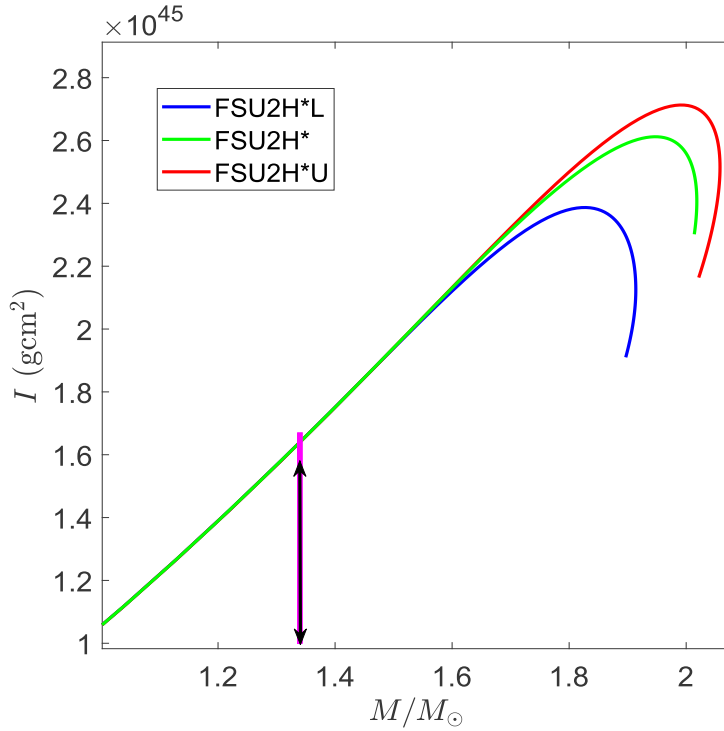


Figure 3.6: Moment of inertia as a function of the star’s mass, obtained with our three parametrizations FSU2H*, FSU2H*L and FSU2H*U. The observational constraints obtained from Ref. [357] are also shown. The figure was published in [309].

inertia, calculated using our three parameterizations, is presented in Fig. 3.6. We compare our results with two sets of constraints outlined in Ref. [357]. It is important to note that these constraints are not direct measurements but are instead obtained utilizing universal relations between NS observables. The constraints shown in the figure employ different ranges of tidal deformability values reported in the more restrictive case (Ref. [352]) and the less restrictive case (Ref. [351]).

The FSU2H* family is not compatible with the tighter constraint. However, our results are nearly consistent, within their respective error bars, with the Bayesian analysis conducted in Ref. [358]. A more reliable constraint on EoS models could come from direct measurements of I . As discussed in the introduction, such measurements may become available in the coming years through radio observations of the double pulsar PSR J0737-3039 [136, 138, 139, 140].

For NSs between $1.0 M_\odot$ and $1.6 M_\odot$, we find that the moment of inertia rises roughly linearly with M . However, this linear trend deviates at higher masses for our softest parametrization, FSU2H*L. Our stiffer FSU2H*U EoS, on the other hand, maintains the tendency for masses $M \lesssim 1.8 M_\odot$. Similar to the case for tidal deformability, this behavior directly reflects the change in the curvature of the $M(R)$ curves. For the most massive cold NSs, we also observe that the predictions from the different parametrizations of the moment of inertia vary $\simeq 10\%$. We also investigate the moment of inertia for PNSs. As

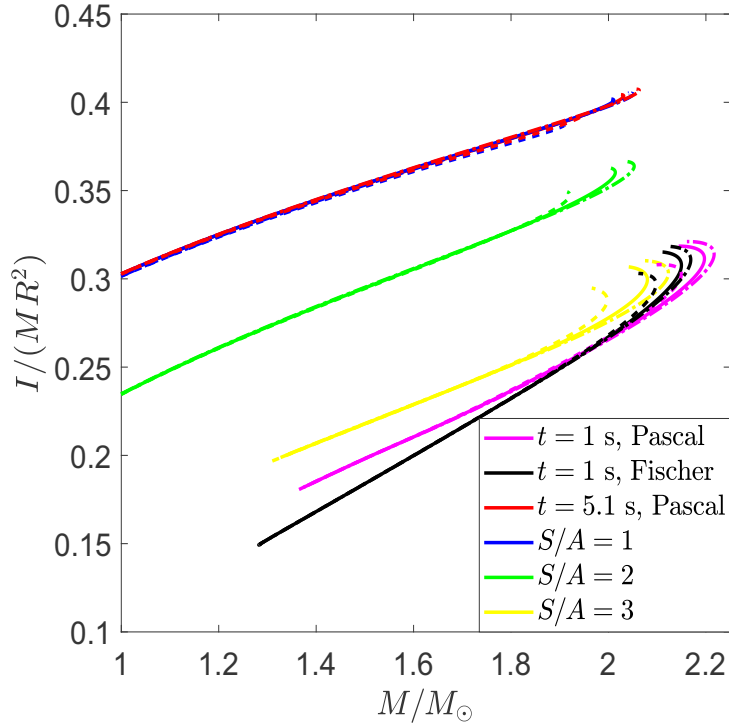


Figure 3.7: Dimensionless moment of inertia obtained with our three parametrizations FSU2H* (solid lines), FSU2H*L (dashed lines), and FSU2H*U (dashed-dotted lines), for PNSs with temperature and lepton fraction as described in the text. The figure was published in [309].

illustrated in Fig. 3.7, a clear trend is observed for the dimensionless moment of inertia, $I/(MR^2)$, as a function of the star's mass: colder stars generally exhibit a higher $I/(MR^2)$.

This is a direct consequence of their significantly smaller radii for a given mass. The dimensionless quantity appears to be less affected by the hyperonic uncertainties, with only minor variations observed between the different parametrizations.

4 Strangeness number of cold and hot NSs

To understand how hyperons and their associated uncertainties affect NS properties, we calculate the strangeness number $\mathbf{S}$ of a star, a quantity that directly relates to the amount of hyperons present in the star. The strangeness number, taking into account the GR metric for a spherical star, is given by:

$$\mathbf{S} = 4\pi \int_0^R \frac{r^2 \sum_i n_i^s}{\sqrt{1 - \frac{2GM(r)}{r}}} dr. \quad (3.12)$$

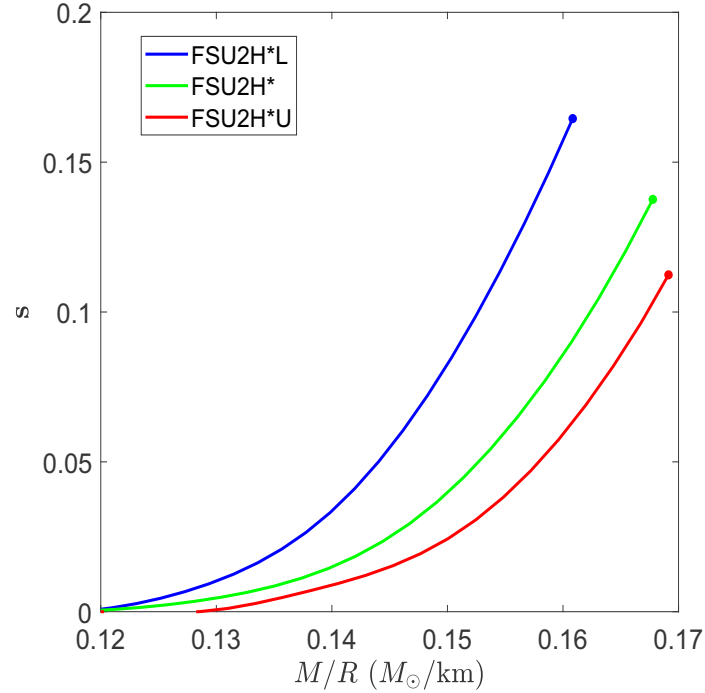


Figure 3.8: Normalized strangeness number s as a function of M/R , computed with our three parametrizations. The dot at the end of each curve indicates the maximum mass star configuration. The graph was published in [309].

Here, $n_i^s = s_i \rho_i$ represents the strangeness density for any baryon i , while s_i is the strangeness quantum number of the baryon. Because $\mathbf{S}$ is a large number, we will show

our results using the normalized strangeness number, denoted as $s = \mathbf{S}/N_B$, where the total baryon number N_B is defined as

$$N_B = 4\pi \int_0^R \frac{r^2 \sum_i \rho_i}{\sqrt{1 - \frac{2GM(r)}{r}}} dr. \quad (3.13)$$

Fig. 3.8 shows the relationship between the normalized strangeness number of a cold star and its compactness. Across all of our model parametrizations, the strangeness number continuously rises with compactness. This clearly illustrates that an increased compactness promotes the creation of hyperons by increasing the star's average and center densities. In addition, softer models lead to a higher strangeness number for a given compactness value. Similar results are obtained for hot stars (Fig. 3.9). However,

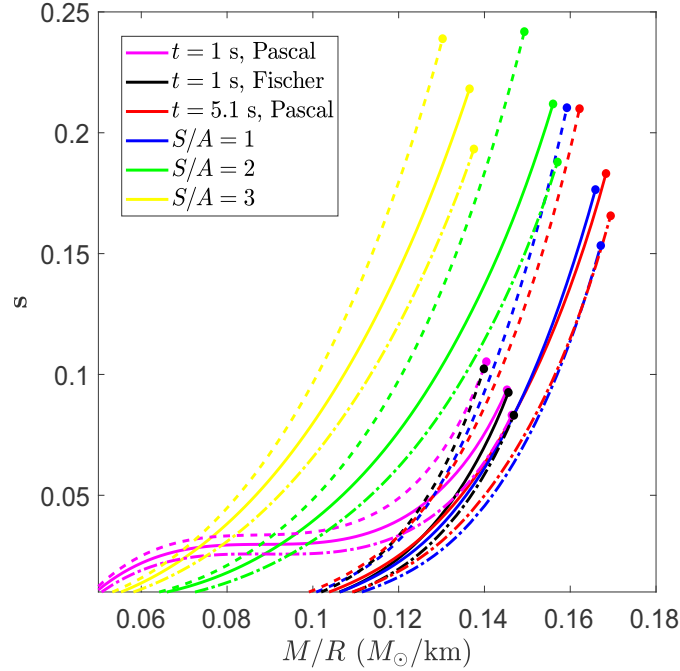


Figure 3.9: Normalized strangeness number s , computed with our three parametrizations: FSU2H* (solid lines), FSU2H*L (dashed lines), and FSU2H*U (dashed-dotted lines), for PNSs with temperature and lepton profiles as described in the text. The dot at the end of each curve indicates the maximum mass star configuration. The graph was published in [309].

the normalized strangeness number is highly sensitive to the temperature profile in the high-density region of the stars. As shown in the figure, the calculations performed at a high constant entropy per particle and with lower lepton fraction ($S/A = 3$), for a given

compactness consistently produce the largest values of strangeness. Conversely, stars with lower temperatures in the high-density region and higher lepton fractions ($t = 1$ s, Pascal and $t = 1$ s, Fischer) require greater compactness to produce hyperonic matter. It is also clear that the normalized strangeness number is influenced by the hyperonic uncertainties, with smaller values for the stiff version of our model and larger values for the soft version.

Chapter 4

Hyperons signatures in binary neutron star (BNS) mergers

In the previous chapters, we presented the FSU2H* model made to describe dense hypernuclear matter at finite temperature. With this model, we analyzed the composition of matter and different thermodynamic quantities. Predictions for several stellar observables at both zero and finite temperature were also presented. We discussed that a large uncertainty still exists in the hyperonic sector of our model parameters, which we attempted to capture using two additional extreme parameterizations. However, the results we presented are limited to our particular model.

As we already discussed in the introduction, many different methods can be used to obtain an EoS of dense and hot matter that is compatible with current constraints. To break this model degeneracy and gain valuable information about the true composition of matter, one can study the observables of NSs. We will now focus on providing a clear pathway for inferring the presence of hyperons in matter using model-independent signatures.

The most obvious observational candidate to start the investigation is the $M(R)$ relation of cold NSs. It is well known that when a hyperonic species appears in matter, the speed of sound, given by

$$c_s^2 = \frac{dP}{d\varepsilon}, \quad (4.1)$$

suffers a drop [359, 360, 361, 362, 363], which is a direct consequence of introducing new particles into the dense medium. The appearance of hyperons leads to a softening of the hyperonic EoS compared to the nucleonic one—that is, an EoS constructed with the same NN interactions but without considering hyperons as a possible degree of freedom. Since a unique correspondence exists between the EoS and the $M(R)$ relation, many hyperonic models consequently fail to reproduce the observed masses of massive pulsars,

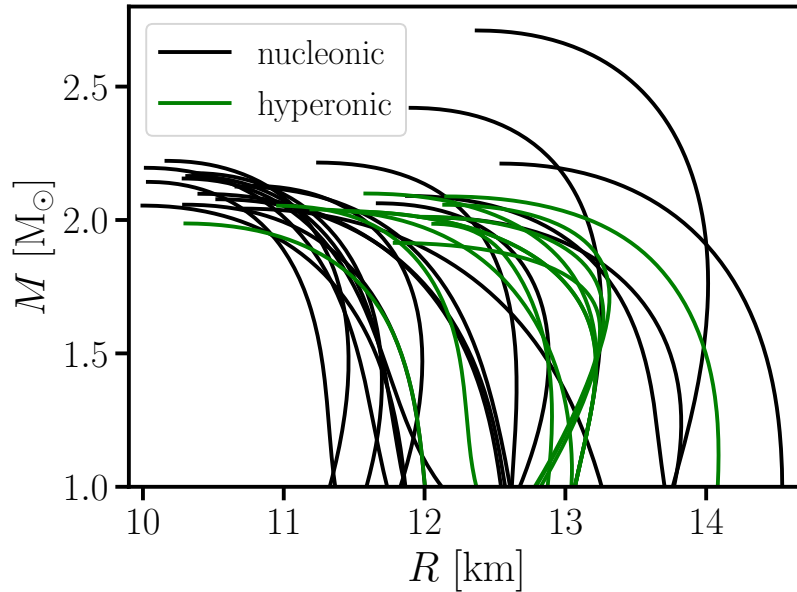


Figure 4.1: Mass-radius relations computed with many different hyperonic (green lines) and nucleonic (black lines) models. The graph was published in [364], and it was made by Sebastian Blacker.

a problem known as the hyperon puzzle [230]. One might therefore conclude that a sufficiently precise measurement of the $M(R)$ relation for cold stars could resolve the question of the composition of dense matter.

However, as shown in Fig. 4.1 from Ref. [364], where the $M(R)$ relation was computed using numerous nucleonic and hyperonic models, the two classes of models do not exhibit distinguishable features. In other words, a similar $M(R)$ relation can be obtained from two models with different degrees of freedom and distinct underlying nuclear interactions. This problem is analogous to the well-known ‘masquerade’ problem in the context of quark matter. As discussed in Ref. [365], EoSs for hybrid stars can mimic the mass-radius relationships of ordinary NSs, making it challenging to extract the composition of matter. Similarly, hyperonic EoSs can mimic nucleonic EoS, implying that even a very precise measurement of the $M(R)$ relation may not yield a definitive answer as to whether strangeness is produced in dense matter¹.

An alternative strategy is to examine the characteristics of hot and dense hypernuclear matter and analyze how the distinctive hyperonic signatures in the thermodynamic

¹While this thesis was being written, a promising way to extract information about the composition of matter based on cold star observables was developed in Ref. [366].

quantities translate to observations where such conditions are met. We have already shown in Chapter 2 that, within the framework of the FSU2H* model, the finite-temperature β -equilibrated EoS of hypernuclear matter exhibits a qualitatively different behavior compared to its nucleonic counterpart. In particular, we demonstrated that the thermal pressure suffers a significant drop when hyperons are present.

In the first part of this chapter we will start by investigating whether these hyperonic features are model-independent and present in matter outside of β -equilibrium. This will be accomplished by studying the thermal properties of 14 distinct hyperonic three-dimensional (3D) EoSs (temperature, baryon density, and charge fraction dependent). Once the signatures are identified, we will analyze how they translate to observables and quantities relevant to BNS merger events. In these events, matter can be heated significantly, making the observables sensitive to the 3D EoS. Therefore, the second part of this chapter will be dedicated to studying the outcome of a simulation campaign of BNS merger events that we have performed using 14 different 3D hyperonic EoSs and 23 different 3D nucleonic EoSs. All simulations that we will present in this chapter were set up by Sebastian Blacker and Georgios Lioutas and were performed using the supercomputers at the GSI theory department in Darmstadt, Germany.

1 Thermal behavior of hyperonic equations of state (EoSs)

This section is dedicated to exploring the thermal properties of the hyperonic EoSs that will be employed in our simulation campaign. The characteristics of the nucleonic models used for comparison have been extensively discussed elsewhere (see, e.g., [226, 228, 229, 243, 244, 245, 367]) and will not be re-examined here. However, the systematic differences in the results obtained using nucleonic/hyperonic models will be frequently highlighted throughout the subsequent discussion.

The EoSs that will be presented are developed in Density Functional Theory (DFT) frameworks, assuming hadron weak equilibrium. The parameters of the models are usually constrained by the nuclear matter and finite nuclei data. The construction makes these models particularly suitable for studying systems close to isospin-symmetric configurations at densities around ρ_0 . They can, however, be utilized to explore denser systems at higher asymmetries and at finite temperatures when they are carefully extrapolated.

1.1 Brief description of hyperonic EoSs

We employ hyperonic EoSs that fall within the broad category of covariant DFT models. These theoretical frameworks are typically expressed in terms of a Lagrangian, which characterizes the interactions between nucleons and hyperons. The complex many-body problem is effectively simplified and is modeled at the mean-field level through the exchange of various virtual mesons.

There are several ways in which one can categorize the different RMF models. Some of them, such as our FSU2H* family and SFHOY [248], incorporate nonlinear self-interaction terms for the scalar and vector mesons (e.g., σ^3 , σ^4 , ω^4). As we have explained in Chapter 2, these terms allow for a more accurate description of nuclear matter properties, particularly its incompressibility and particles' effective masses. Furthermore, they can influence the stiffness of the EoS at high densities. In contrast, other models, including DD2Y [273], BHBA ϕ [272], the DDLSY family [258], and DD2Y Δ [251], utilize density-dependent couplings. In these models, the coupling strengths between hadrons and mesons are not constant but rather depend explicitly on the baryon density. This density dependence often mimics the effects of higher-order interactions or medium modifications, offering an alternative approach to control the high-density behavior of the EoS and to provide a more flexible description of finite nuclei and nuclear matter properties.

Beyond the general classification of RMF models, different groups of these EoSs are designed to explore the uncertainties in baryon-baryon interactions. We have already shown that the FSU2H* family was created for a comprehensive exploration of the uncertainties inherent in the YN and YY interactions. The DDLSY family, on the other hand, is designed to investigate the uncertainty associated with the slope of the symmetry energy at high densities. The three different parametrizations, DDLSY(30), DDLSY(50), and DDLSY(70), differ by the predicted slope parameter. Furthermore, the DD2Y Δ family expands the hadronic degrees of freedom by including Δ resonances, in addition to the conventional baryon octet members. These models are constructed under two simplifying assumptions. First, it is assumed that the vacuum masses of the Δ resonances are not significantly altered by hadronic interactions in a dense medium. Second, Δ s are assumed to have negligible widths. It will be demonstrated in the subsequent discussions that the appearance of Δ resonances at high densities has an effect on the EoS equivalent to that produced by hyperons, both at zero and at finite temperature. For this reason, we have chosen to include them in our work.²

²Please note that whenever we refer to 'hyperonic' models, we are referring to the models that consider not only hyperons as an additional hadronic degree of freedom, but also to the models that also account for Δ resonances.

As a distinct type of covariant DFT, we also consider the SU(3) Chiral Mean Field (CMF) model, more precisely, the DNS EoS [368]. Unlike other RMF approaches, the CMF model is rooted in the fundamental symmetries of QCD, particularly its chiral symmetry. In vacuum and at low densities, this symmetry is spontaneously broken, which generates the large masses of the baryon octet. However, an important feature of the CMF model, relevant for the extreme conditions inside NS matter, is its prediction of chiral symmetry restoration at sufficiently high temperature and/or density. As matter becomes hotter and denser, the chiral condensate, which acts as the order parameter for chiral symmetry breaking, melts. This leads to a reduction in the effective masses of the baryons, which tend towards their smaller 'chiral limit' or 'bare' masses. The process is often modeled as a smooth crossover at finite temperature but could potentially be a first-order phase transition at very high baryon densities and low temperatures.

We also incorporate the Quark-Meson-Coupling (QMC) model, specifically the QMC-A parametrization [256], into our study. This DFT model represents a distinct subclass of the covariant models, as it explicitly incorporates the internal quark structure of hadrons to describe their dynamics. While other RMF models treat hadrons as point-like particles, the QMC model postulates that hadron interactions in baryon matter arise from the direct coupling of their constituent valence quarks to meson fields. In this framework, nucleons and hyperons are considered as bags of quarks confined by the strong force. The coupling of quarks to mesons effectively modifies the properties of the hadrons within the dense medium. This approach offers several advantages. It inherently links hadron properties to the quark-quark interactions and provides a more consistent description of the in-medium hadron modifications, such as changes in their sizes and effective masses. It also naturally incorporates aspects of confinement and allows for a more direct connection to QCD at intermediate densities.

In Fig. 4.2 we present the onset density of heavy baryons as a function of temperature for all considered models, assuming matter in β -equilibrium. The plot allows us to precisely quantify the conditions under which heavy baryons begin to play a significant role in dense matter at finite temperature. While a non-zero, albeit minuscule, fraction of heavy baryons will always be present due to thermal excitations at any finite temperature, we define the onset density as the value at which the amount of heavy baryons in matter exceeds 1% of the total baryon density.

At low temperature, relevant for the initial inspiral phase of a binary merger, most models predict hyperon appearance around $2\rho_0$. This suggests that dense NS cores of more massive stars (usually with masses $M \gtrsim 1.5M_\odot$) are likely to contain hyperons. However, for some of the models that incorporate Δ resonances, the onset density can be as low as $\lesssim 1.5\rho_0$. This early appearance is driven by the energetically favorable formation of negatively charged Δ^- resonances, which, similar to electrons, can contribute to maintaining charge neutrality at lower baryon chemical potentials compared to other

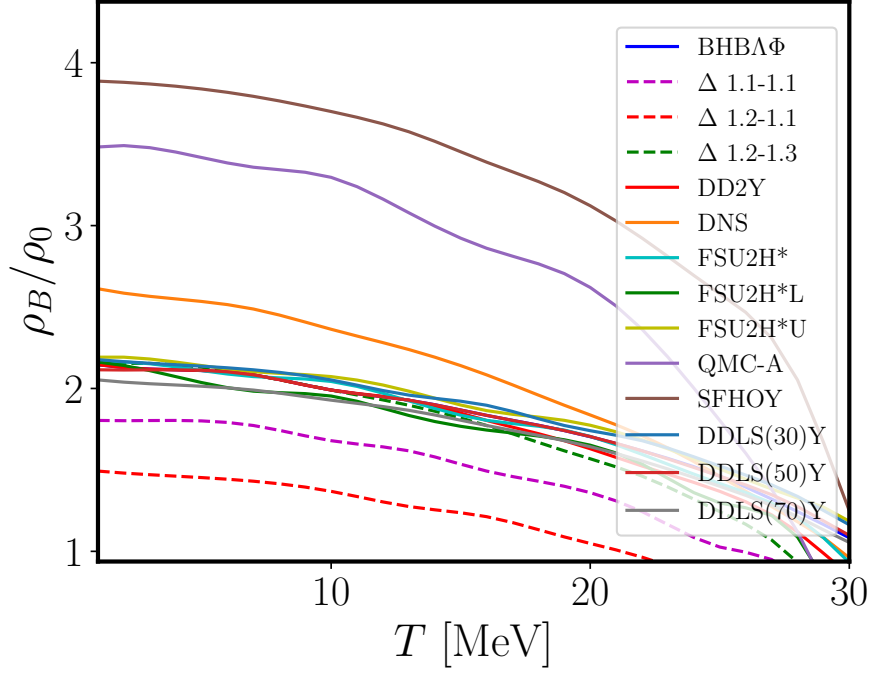


Figure 4.2: Onset density of heavy baryons as a function of temperature, predicted by the different models utilized in the study. The onset density is defined as the density at which the abundance of exotic species in matter is larger than 1%. The figure was published in [369].

hyperons. The early presence of Δ^- baryons, therefore, introduces a softening mechanism to the EoS at densities below the typical hyperon onset, leaving open the possibility of heavy baryon signatures in inspiral observables, even from low-mass NSs. In contrast, certain models, such as SFHOY and QMC-A, exhibit a delayed onset, pushing hyperon appearance to densities larger than $3\rho_0$. For those EoSs, the inspiral observables are likely insensitive to the appearance of hyperons and Δ resonances.

A universally observed trend across all models is that increasing temperature consistently shifts the onset density to lower values. Of course, this is a direct consequence of thermal excitations, which provide sufficient energy to overcome the mass difference between nucleons and hyperons, thereby enabling their population at lower densities in contrast to the zero temperature case. At a moderate temperature of $T = 30$ MeV, typically reached in the post-merger remnant of BNS mergers, we observe that the onset density is reduced to values lower than $1.5\rho_0$ for all models.

1.2 Density dependence of thermodynamic quantities

In this subsection we will investigate the density dependence of several thermodynamic quantities at finite temperature for all the previously discussed hyperonic models. We will focus on matter at a fixed temperature $T = 25$ MeV and charge fraction $Y_Q = 0.1$. In addition, for comparison, the results of three nucleonic models will be shown in the analysis. These include SFHO [271, 370] (represented by a solid black line), DD2 [338, 271] (dashed black line), and FSU2R [259] (dashed-dotted black line). The selected nucleonic models are also constructed in an RMF framework, and they share many properties with some of the specific families of the hyperonic EoSs we employ. For instance, the DD2 model serves as a direct nucleonic baseline for the models BHB $\Lambda\phi$, DD2Y, and the DD2Y Δ family. Similarly, SFHO exhibits common features with SFHOY, while FSU2R is closely related to FSU2H*. Both FSU2R and the FSU2H* family belong to the group of RMF models with non-linear self-interactions and have their parameters constrained in an analogous way.

We will begin by analyzing the thermal contribution to the energy density (ϵ_{th}) and the pressure (P_{th}). For any quantity X , the thermal contribution can be defined as the difference between the value of the quantity at finite temperature T , density ρ_B , and charge fraction Y_Q and the value of the same quantity under equal conditions of ρ_B and Y_Q but at $T = 0$:

$$X_{\text{th}} = X(T, \rho_B, Y_Q) - X(0, \rho_B, Y_Q). \quad (4.2)$$

Fig. 4.3 illustrates the thermal energy density ϵ_{th} as a function of the baryon density. All hyperonic EoSs, except QMC-A, predict that ϵ_{th} is a monotonously increasing function. This general trend is physically intuitive: as the density of matter increases, there are more particles per unit volume available to contribute to the thermal energy through their kinetic motion and excitation into higher energy states. It is noticeable that models incorporating a larger number of heavy baryon species (i.e., those with hyperons and Δ resonances appearing at lower densities) systematically exhibit higher values of the thermal energy density. The additional degrees of freedom provide more channels for thermal excitation and energy storage. When hyperons or Δ baryons become energetically accessible, their thermal populations contribute significantly to the total energy density, as they are less degenerate than the nucleons. This in turn leads to a substantial increase in ϵ_{th} .

The behavior of P_{th} is much more interesting than the one of ϵ_{th} . As seen in Fig. 4.4, a counterintuitive effect is observed. Contrary to the purely nucleonic models, which generally exhibit a relatively featureless behavior, the hyperonic models show a significant and striking reduction of $P_{\text{th}}(\rho_B)$. This reduction manifests as a distinctive dip, or a series of dips, occurring precisely at the onset density for each new species at zero

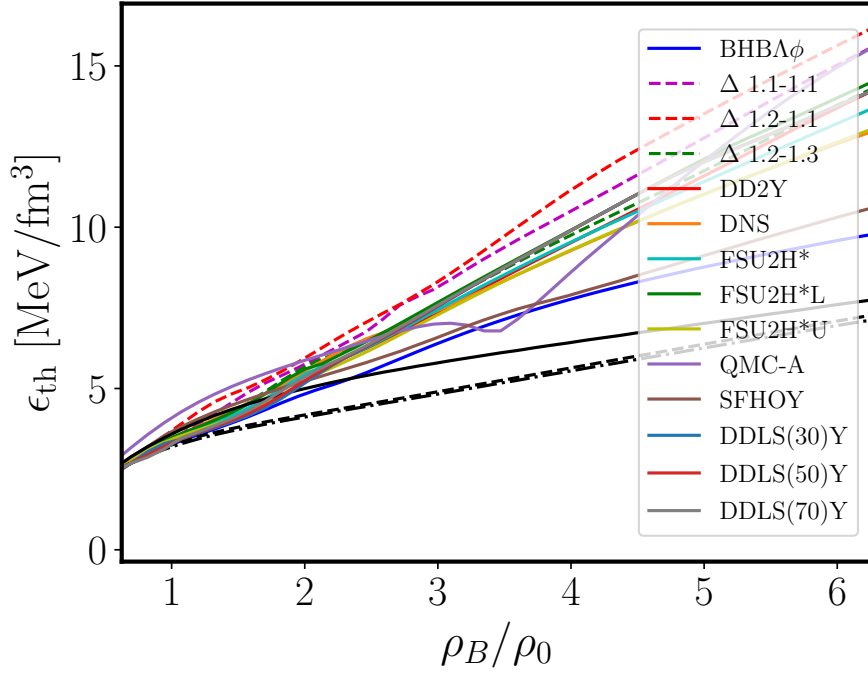


Figure 4.3: Thermal energy density ϵ_{th} as a function of density at fixed charge fraction $Y_Q = 0.1$ and temperature $T = 25$ MeV. Three different nucleonic models (solid line for SFHO, dashed line for FSU2R, and dash-dotted line for DD2) are also shown. The figure was published in [369].

temperature. We have already discussed this interesting feature in Chapter 2, when β -stable matter was studied within the framework of the FSU2H* model. These results have been also shown in some recent works [251, 289]. It is fundamentally related to the redistribution of particle populations and, consequently, to the reduction of degeneracy pressure that occurs when new species are introduced in matter. As Figure 4.2 has already demonstrated, at finite temperature, the available thermal energy makes it energetically favorable for heavier baryons to appear at substantially lower densities than at zero temperature. Those hyperons appear at the cost of the most energetically present nucleons in nucleonic matter. While the newly populated heavy baryons contribute their own thermal kinetic pressure, this contribution is often outweighed by the significant decrease in the degeneracy pressure of the previously dominant light baryons, resulting in an overall net dip in the total thermal pressure. Models like DD2Y Δ 1.1 – 1.1, DD2Y Δ 1.2 – 1.1, or those of the DDLSY family, which predict the appearance of different heavy baryon species at distinct densities, show a sequence of such dips. Each dip corresponds to the successive onset of a new species. This complex behavior of the thermal pressure leads to a softening of the finite-temperature EoS in the presence of exotic degrees of

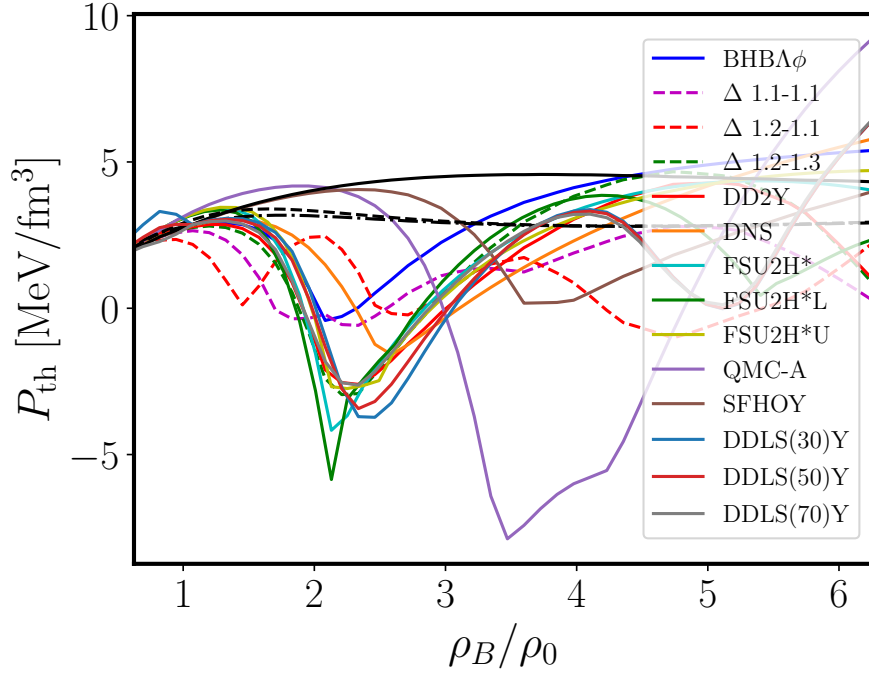


Figure 4.4: Thermal pressure P_{th} as a function of density at fixed charge fraction $Y_Q = 0.1$ and temperature $T = 25$ MeV. Three different nucleonic models (solid line for SFHO, dashed line for FSU2R, and dash-dotted line for DD2) are also shown. The figure was published in [369].

freedom. We again reiterate that hadrons are considered to be in weak equilibrium. This implies that for a fixed charge fraction Y_Q , the equilibrium composition of matter is highly sensitive to temperature.

To quantify and analyze the different thermal behavior of the thermal energy density and thermal pressure we introduce the thermal index, $\Gamma_{\text{th}}(\rho_B, Y_Q, T)$, for 3D EoSs, which is defined in an analogous way as in Chapter 2:

$$\Gamma_{\text{th}}(\rho_B, Y_Q, T) = 1 + \frac{P_{\text{th}}(\rho_B, Y_Q, T)}{\varepsilon_{\text{th}}(\rho_B, Y_Q, T)}. \quad (4.3)$$

In the context of BNS mergers, the thermal index can be seen as a quantity that measures how efficiently the thermal energy stored in the system is converted into thermal pressure. The results obtained with our set of EoSs are given in Fig. 4.5. The thermal index of purely nucleonic models exhibits only a soft density dependence, typically staying relatively close to values between 1.5 and 2. On the contrary, the hyperonic Γ_{th} behavior is profoundly different. The previously discussed behavior of the thermal pressure directly

translates into an abrupt drop in the thermal index. It is noteworthy that, for most of the models including exotic degrees of freedom, the thermal index falls to values lower than 1.5 in the intermediate- and high-density regions. Since this feature is model-independent, it can have significant implications for numerical simulations of BNS mergers. As we have already discussed, many astrophysical simulations often employ a simplified 'ideal gas' approximation for the thermal part of the EoS, assuming a constant thermal index typically around $\Gamma_{\text{th}} \approx 1.8 - 2$. Our results clearly demonstrate that for hypernuclear matter this approximation can be inaccurate. A thermal index significantly lower than 1.5 suggests that the thermal stiffening of the EoS is less pronounced than commonly assumed, leading to a 'softer' thermal response.

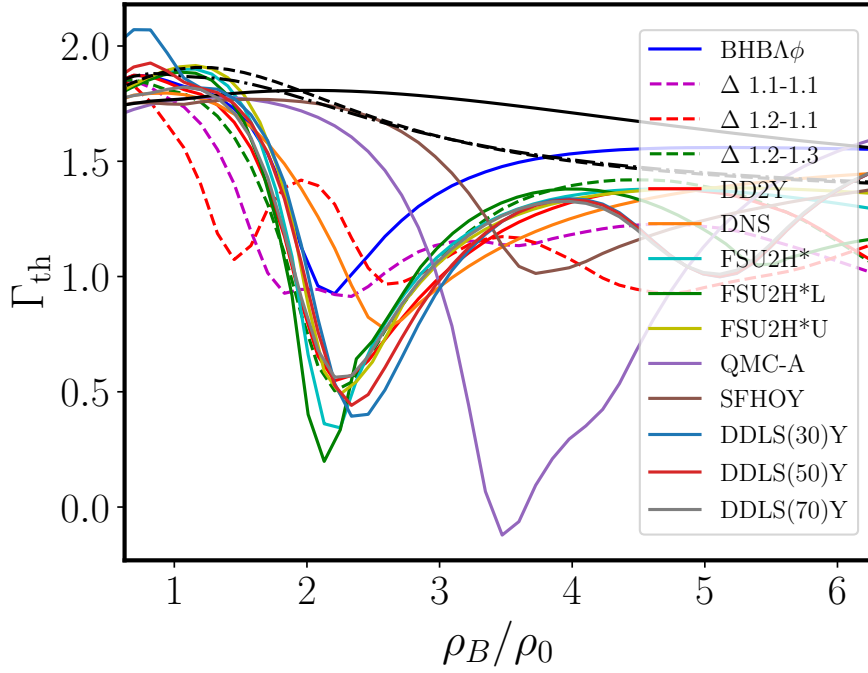


Figure 4.5: Thermal index Γ_{th} as a function of density at fixed charge fraction $Y_Q = 0.1$ and temperature $T = 25$ MeV. Three different nucleonic models (solid line for SFHO, dashed line for FSU2R, and dash-dotted line for DD2) are also shown. The figure was published in [369].

To further confirm our conclusion that the observed reduction in thermal pressure and the characteristic dip in the thermal index are fundamentally attributed to the loss of degeneracy pressure, we introduce and analyze the normalized hyperon excess, $\Delta\rho_Y$. The hyperon excess quantifies the thermal population of heavy baryons, and it can be

computed using the following relation:

$$\Delta\rho_Y = \frac{\sum_{Y_i}(\rho_{Y_i}(T, \rho_B, Y_Q) - \rho_{Y_i}(0, \rho_B, Y_Q))}{\rho_B}. \quad (4.4)$$

In the last equation, ρ_{Y_i} denotes the density of an arbitrary heavy baryon species Y_i . This quantity effectively represents the fraction of baryons that exist purely due to thermal excitations and the opening of new degrees of freedom in matter. The hyperon excess, computed for all hyperonic models at our disposal, is displayed in Fig. 4.6. A

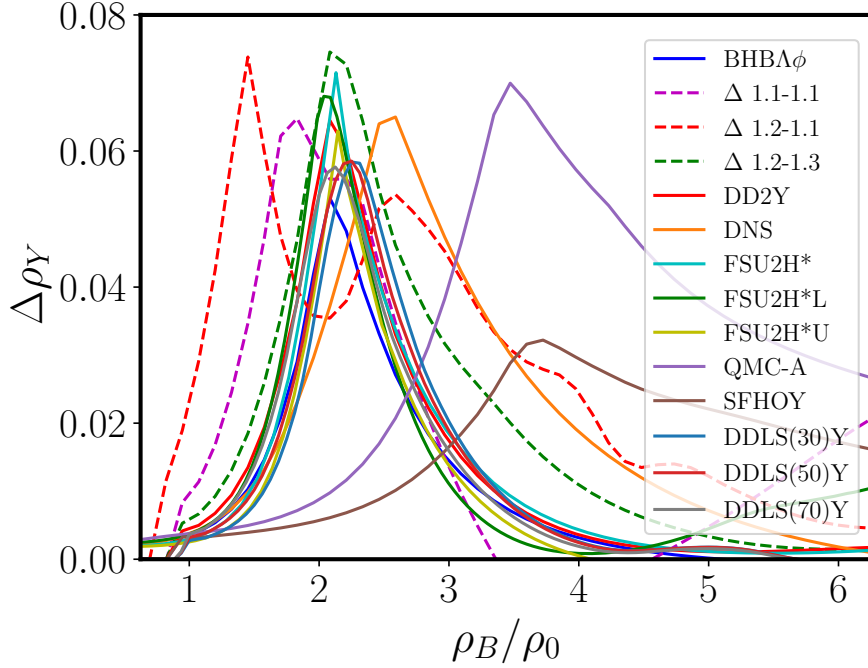


Figure 4.6: Hyperon excess $\Delta\rho_Y$ as a function of density at fixed charge fraction $Y_Q = 0.1$ and temperature $T = 25$ MeV. The figure was published in [369].

compelling correlation emerges when one compares the thermal index with the hyperon excess. Indeed, the densities where the hyperon excess exhibits a distinct peak occur at the same densities where the prominent dips in the thermal index are observed. The formation of 'excess' hyperons at finite temperature, which are otherwise suppressed at zero temperature, comes at the energetic cost of reducing the highly degenerate populations of nucleons. As the hyperon excess begins to decrease beyond its peak, the hyperons start to become more abundant even at $T = 0$, the influence of the degeneracy pressure redistribution diminishes, and the thermal index consequently increases again, approaching values more typical of nucleonic matter. The exact shape of the hyperon excess as a function of density is highly model-dependent. Some models, like BHBA ϕ

and FSU2H* display relatively sharp peaks. This indicates that they tend to lose their excess amount of thermal hyperons with increasing density faster, leading to a rapid re-stiffening of their thermal index. Conversely, models such as SFHOY and QMC-A have wider peaks, indicating that they maintain a substantial thermal hyperon excess over a broader range of densities. This results in lower values of Γ_{th} persisting across a larger density interval.

The behavior of the DD2Y Δ 1.1 – 1.1 model is particularly insightful, as it demonstrates the effects of multiple exotic species with considerably different onset densities. Indeed, three sequential peaks in $\Delta\rho_Y$, corresponding to the successive onset of different Δ baryons and hyperons, can be observed. The largest peak in $\Delta\rho_Y$ might appear at a relatively low density. Nevertheless, the reduction of the thermal index is more pronounced at higher densities. This rather odd result can be understood by considering that the degeneracy pressure of any species depends strongly on its particle density (approximately as $\rho^{5/3}$ for a non-relativistic, non-interacting degenerate gas). Therefore, the magnitude of the reduction in degeneracy pressure because of the population shifts is more significant when occurring at higher densities, even if the relative excess might be smaller for species with higher onset density.

We finish our discussion on the thermal behavior of hyperonic EoSs by investigating the behavior of the specific heat at constant pressure c_p , as presented in Fig. 4.7. This quantity is related to the temperature derivative of the entropy per particle (S/A) at constant pressure, baryon number, and charge fraction:

$$c_p = T \left. \frac{d(S/A)}{dT} \right|_{P,A,Y_Q}. \quad (4.5)$$

Physically, c_p quantifies the amount of heat energy per particle required to raise the system's temperature by 1 MeV while maintaining the pressure constant. A higher specific heat implies that a system can absorb more thermal energy for a smaller increase in temperature. From the figure, it is evident that in the high-density region, hyperonic EoSs consistently produce higher values of c_p compared to their purely nucleonic counterparts. This difference is related to the increased capacity of hypernuclear matter to store thermal energy due to the generation of new degrees of freedom [310, 251, 289].

In summary, our investigation of the thermal behavior of hyperonic EoSs has revealed distinct features, such as reduced (and in some cases even negative!) thermal pressure and higher specific heat at constant pressure in matter that significantly differentiate them from purely nucleonic models. In the subsequent section, we will explore whether these unique microphysical imprints leave signatures on the macroscopic observables and key quantities relevant to BNS merger events.

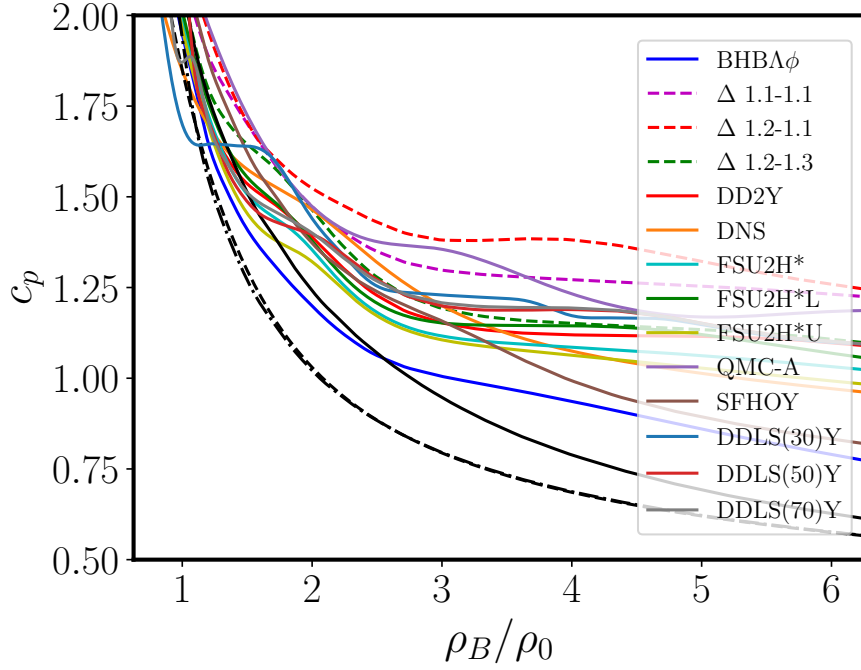


Figure 4.7: Specific heat at constant pressure c_p as a function of density at temperature $T = 25$ MeV and fixed charge fraction $Y_Q = 0.1$. Three different nucleonic models (solid line for SFHO, dashed line for FSU2R, and dash-dotted line for DD2) are also shown. The figure was published in [369].

2 BNS merger simulations

The specific thermal features of hyperonic EoSs that we have identified in the previous section can leave imprints on observables and quantities that are sensitive to the finite-temperature EoS. To test this possibility, we have run a large simulation campaign. The setup and the results of the campaign are explained in this section.

2.1 Approach

Most simulations presented in the thesis are performed for symmetric binaries, namely $q = M_1/M_2 = 1$ with a total mass of $M_{\text{tot}} = 2.8 M_{\odot}$. In some cases, specifically highlighted, we will present results from simulations of asymmetric binaries for stars with $q = M_1/M_2 = 0.8$; $M_{\text{tot}} = 2.8 M_{\odot}$, and also more massive symmetric binaries. As a starting point for each simulation, we choose the configuration where the centers of the stars are separated by 38 km. The stars are moving on quasi-equilibrium circular

orbits and are not rotating around their axes. Initially, the stellar matter is cold and in β equilibrium.

To solve the coupled hydrodynamic and field equations, the Smoothed Particle Hydrodynamics (SPH) code is used that has been described in detail in Refs. [371, 372, 373]. Each simulation is performed with 300,000 SPH particles. The field equations are solved in the so-called Conformal Flatness Approximation (CFA) [374, 375]. The CFA assumes that the spatial part of the three-dimensional metric can be written as a conformally flat metric multiplied by a conformal factor. The full Einstein equations are replaced with a system of elliptic equations for the lapse function, shift vector, and conformal factor. To account for the GW emission in CFA, a back-reaction formalism is implemented that mimics the energy and angular momentum dissipation by GWs [376, 377]. We also assume the neutrino emission to be negligible [372].

All of the EoSs that we have at our disposal are 3D and roughly compatible with all known astrophysical constraints. To properly discriminate the effects that hyperons have on certain thermodynamic quantities and astrophysical observables, simulations with two large sets of EoS are performed. The first sample, which we will refer to as hyperonic, consists of the 14 different EoS that we have presented at the beginning of this section: BHBA ϕ [272], DD2Y [273], DDLSY family [258], DNS [368], FSU2H* family [289, 309], QMC-A [256], DD2Y Δ family [251], and SFHOY [248]. The baryons are assumed to be in weak equilibrium for all EoS considered in our hyperonic sample. The second sample, which we will refer to as nucleonic, consists of 23 different EoSs that do not consider exotic components. The EoSs considered in this set are the following ones: APR [378, 379], DD2 [338, 271], DD2F [338, 380], DDLSN family [258], FSU2R [259], FTNS [333, 381], GS2 [336], LPB [382, 383], LS220 [156], LS375 [156], SFHO [271, 370], SFHX [271, 370], SRO(SLy4) [384, 385], TM1 [386, 387], TMA [388, 387], and the models 'Fiducial', 'Large Mmax', 'Large SL', 'Large R', 'Small SL', and 'Smaller R' from Ref. [389].³ We note that of all EoSs in our hyperonic set, only DNS and BHBA ϕ are compatible with the mass $M > 2.09 M_{\odot}$ (3σ confidence level) of the massive black widow pulsar PSR J0952-0607. [327]. Similarly, the DD2F, FSU2R, LS220, SFHO, SRO(SLy4) and TMA EoSs from our nucleonic set cannot satisfy this strong constraint.

As previously discussed, the finite-temperature behavior of hyperonic EoSs is qualitatively different from that of the nucleonic ones. To properly analyze how this different temperature behavior affects the different observables and quantities in a BNS merger event, with each EoS of both samples, two runs of simulations were performed. In the first run of simulations, which we will refer to as a 3D run, the full 3D EoS of each set was used. These are simulations where the finite-temperature behavior of the EoS is treated

³The majority of the EoSs that we use can be downloaded freely through the CompOSE website (<https://compose.obspm.fr/>) [244, 341, 342, 390].

consistently, as predicted by the microphysical model itself. Since neutrino transport is not considered, the electron fraction of each fluid element is determined by the initial conditions and is advected with the moving fluid. In the second run of simulations, which we will refer to as a $\Gamma_{\text{th}} = 1.75$ run, the pressure and the energy density contributions are separated into cold and thermal components. That allows for use of Eq. (2.40) to obtain a new set of thermal EoSs. The final EoSs that are used in this run (the $P(\varepsilon)$ relation) are derived assuming cold, β -equilibrated matter. The finite-temperature component is then approximated using a constant thermal index of $\Gamma_{\text{th}} = 1.75$ in Eq. (2.40). In this run, the lepton fraction is always set to the cold beta-equilibrium value at the respective density (see Ref. [316] for more details).

As we will show *a posteriori*, this simplified approach accurately mimics the behavior of the 3D nucleonic EoSs. In that sense, the $\Gamma_{\text{th}} = 1.75$ run serves as a reference run that will be used when the observables from BNS mergers are studied (Section 3). This run shows what the simulation output would look like if the thermal behavior of the corresponding EoS were "ideally nucleonic." The difference in any observable obtained with the 3D run and the $\Gamma_{\text{th}} = 1.75$ run can be attributed to the specific thermal behavior of the corresponding EoS. In an optimistic future scenario, where the cold part of the EoS is known with high accuracy (for instance, from precise $M(R)$ measurements of many pulsars), this difference could be exploited to break the composition degeneracy that exists when only the cold EoS is considered. If certain BNS observables can be measured with high precision, then the actual measurements could be compared with the reference, $\Gamma_{\text{th}} = 1.75$ simulation results to infer the presence of hyperons in dense matter.

2.2 Maximum density, hyperon fraction, thermal index and temperature evolution in BNS merger simulations

In this subsection, we present the results of the time evolution of several quantities, such as the maximum density, hyperon fraction, thermal index and temperature during the merger of two $1.4M_{\odot} - 1.4M_{\odot}$ NSs. Our focus is on the specific behavior of these quantities when hyperons are present, that is, when using hyperonic EoSs, compared to the outcome when we consider three exemplary nucleonic models (DD2, FSU2R, and SFHO). We will present the results obtained from the 3D run only.

We start by looking at the time evolution of the maximum rest-mass density within the merger remnant, $\rho_{\text{max}}(t)$ (Fig. 4.8). The instant $t = 0$ represents the time of merging of the two stars, which in the simulations is characterized by a maximum amplitude in the quadrupole deformation. We can see that the maximum density both during the inspiral phase and throughout the post-merger dynamics is strongly model-dependent. There is a correlation between the EoS stiffness and the initial post-merger dynamics. Softer EoS models, which can be identified by the high $\rho_{\text{max}}(t)$ that they predict already during

the late inspiral phase (SFHOY and DD2Y Δ 1.1 – 1.1), exhibit particularly strong and long-lasting quasi-radial oscillations in the merger remnant. This behavior is consistent with prior studies that have linked EoS softness to the excitation of these oscillation modes [391]. This effect, however, cannot be directly related to the appearance of hyperons. For example, the SFHO model, the nucleonic counterpart of SFHOY, also exhibits strong oscillations. The quasi-radial oscillations of this model are, in fact, stronger than those observed in several other stiffer hyperonic EoSs (the FSU2H* family, for example). We conclude that although hyperons can contribute to the softening of the EoS, an intrinsically soft nucleonic EoS can still lead to more pronounced oscillations than a stiff hyperonic one. Therefore, the effect of hyperons on $\rho_{\max}(t)$ must be interpreted within the context of the total EoS stiffness, emphasizing that hyperon production, by softening the EoS, generically leads to more pronounced oscillations when compared to a nucleonic EoS of similar stiffness at lower densities.

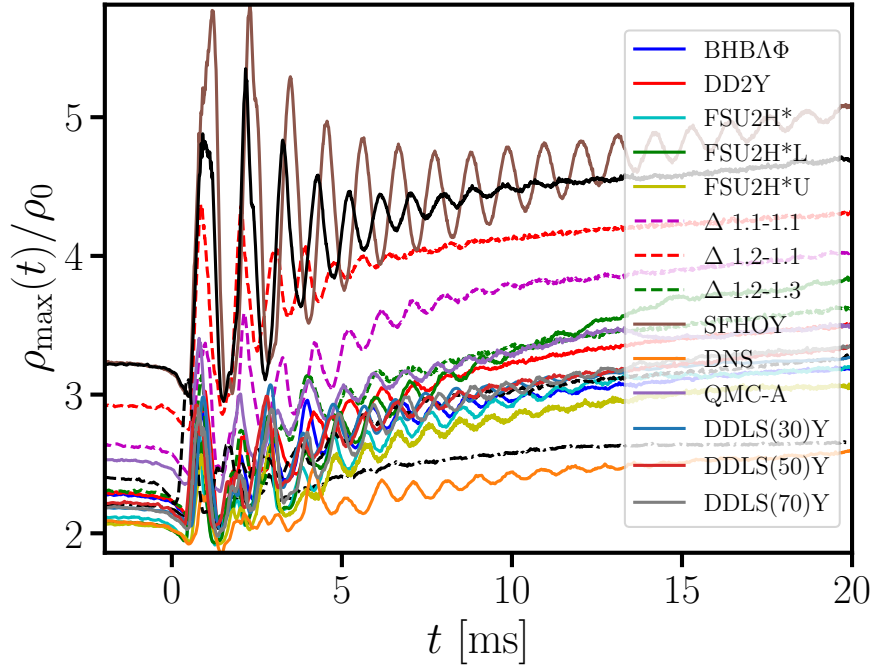


Figure 4.8: Time evolution of the maximum rest mass density in units of nuclear saturation density, $\rho_{\max}(t)/\rho_0$ for different hyperonic EoSs. Three different nucleonic models (solid line for SFHO, dashed line for FSU2R, and dash-dotted line for DD2) are also shown. The figure was published in [369].

Next, we discuss the evolution of the averaged hyperon fraction, which is computed by the following expression:

$$Y_{\text{av}} = \frac{\sum_i m_i Y_i}{\sum_i m_i}, \quad (4.6)$$

where m_i and Y_i are the mass and the hyperon fraction, respectively, of the i -th SPH particle, and the sum is over each SPH particle in the remnant. The results are presented in Fig. 4.9. The behavior of the average hyperon fraction follows the one of the maximum rest-mass density. Of course, this is trivially connected to the dependence of the hyperon onset on the density, as larger maximum densities facilitate greater hyperon production. The analysis of the pre-merger phase reveals that only two of our considered models,

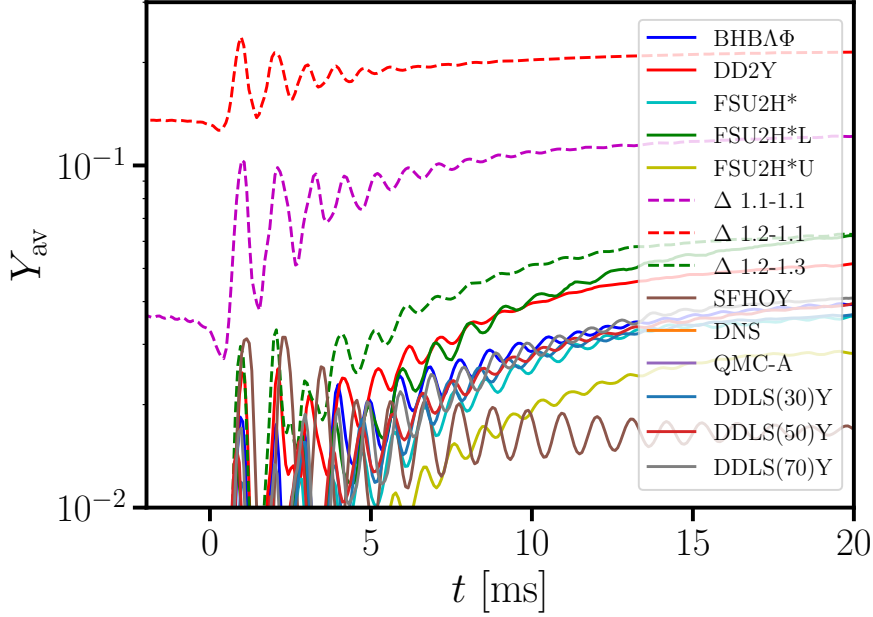


Figure 4.9: Time evolution of the fraction of hyperons in the remnant for different hyperonic EoSs. Please note the logarithmic scale on the y-axis. The figure was published in [369].

namely those from the DD2Y Δ family (DD2Y Δ 1.1 – 1.1 and DD2Y Δ 1.2 – 1.1), predict the presence of exotic species in significant quantities within the NS matter prior to merging. However, for the vast majority of other hyperonic models, the central densities of the cold stars in the inspiral are simply not sufficiently high to induce a population of (non-thermal) hyperons. Hence, the observables that are produced solely from the pre-merging phase, such as the inspiral GW signal, may not carry any information about the exotic baryons

The situation changes in the post-merger phase. As the binary system coalesces and the hyper-massive NS forms, hyperons are predicted to be produced by most of the models. The only exceptions are the QMC-Q and DNS models, which consistently show negligible hyperon fractions. As we will see in the next section, the behaviour of these

two hyperonic models will be similar to the one of the nucleonic schemes during the BNS merger.

Next, in Fig. 4.10, we show the time evolution of the mass-averaged thermal index. This quantity is defined analogously to Y_{av} :

$$\Gamma_{\text{th}}^{\text{av}} = \frac{\sum_i m_i \Gamma_{\text{th},i}}{\sum_i m_i}. \quad (4.7)$$

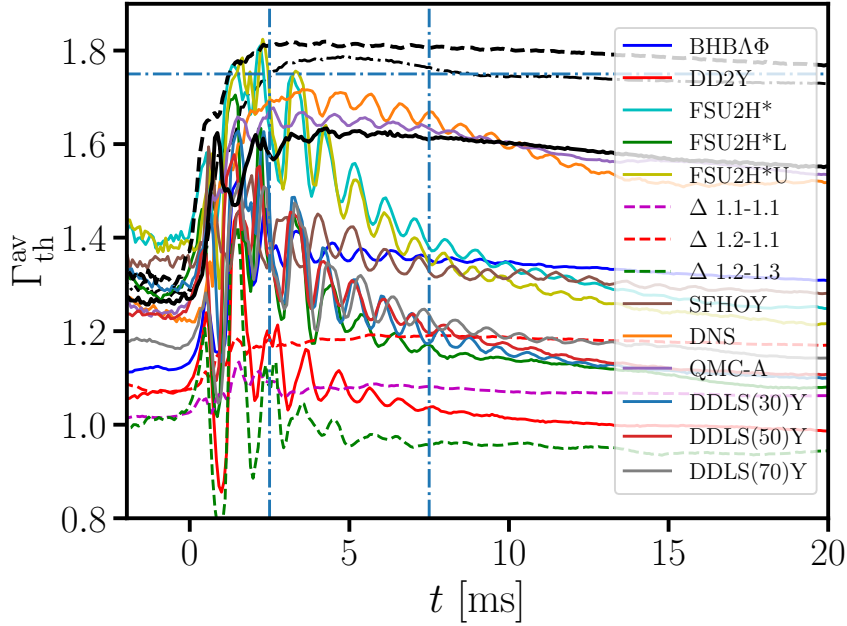


Figure 4.10: Time evolution of the average thermal index for different hyperonic EoSs. Three different nucleonic models (solid line for SFHO, dashed line for FSU2R, and dash-dotted line for DD2) are also shown. The value of $\Gamma_{\text{av}}^{\text{th}} = 1.75$ is marked with a horizontal blue line. The figure was published in [369].

In the inspiral phase, the value of $\Gamma_{\text{th}}^{\text{av}}$ comes from the numerical heating in the simulations, and it does not carry any significance. In the early post-merger phase, for most of our hyperonic EoSs, $\Gamma_{\text{th}}^{\text{av}}$ is characterized by pronounced, large-amplitude oscillations that are correlated with those in the hyperon fraction (cf. Fig. 4.9). The oscillations of the thermal index directly reflect the dynamic interplay between the heating of the remnant, the subsequent thermal population of heavy baryons, and the resulting redistribution of the degeneracy pressure within matter. As new species appear and disappear with the changing density during oscillations, the efficiency with which the thermal energy is converted into pressure support fluctuates, leading to the observed variations in $\Gamma_{\text{th}}^{\text{av}}$.

At later times, when the hyperon fraction fluctuations disappear, the thermal index of the hyperonic EoSs also stops oscillating. On the other hand, nucleonic models either do not predict oscillations or they are significantly smaller in amplitude compared to their hyperonic counterparts. As anticipated from our earlier discussion of the thermal EoS properties (Fig. 4.5), the lowest value of the averaged thermal index is predicted by models with a significant amount of hyperons in the remnant. In fact, three of our models (DD2Y Δ 1.2 – 1.3, DD2Y Δ 1.1 – 1.1 and DD2Y) predict $\Gamma_{\text{th}}^{\text{av}} < 1.2$ just three seconds after the merging. Even for the rest of the models, the value of the average thermal index in the remnant is consistently lower than 1.4 in the later phases of the evolution. The only two models that show a rather high thermal index are DNS and QMC-A, the same two models that did not predict the appearance of hyperons in significant quantities in this merger configuration.

The implications of such low thermal index values could be important in the context of the stability of the remnant. A low value of the thermal index implies that the thermal pressure contribution is significantly reduced, making the overall finite-temperature EoS softer and the remnant more prone to black-hole collapse.

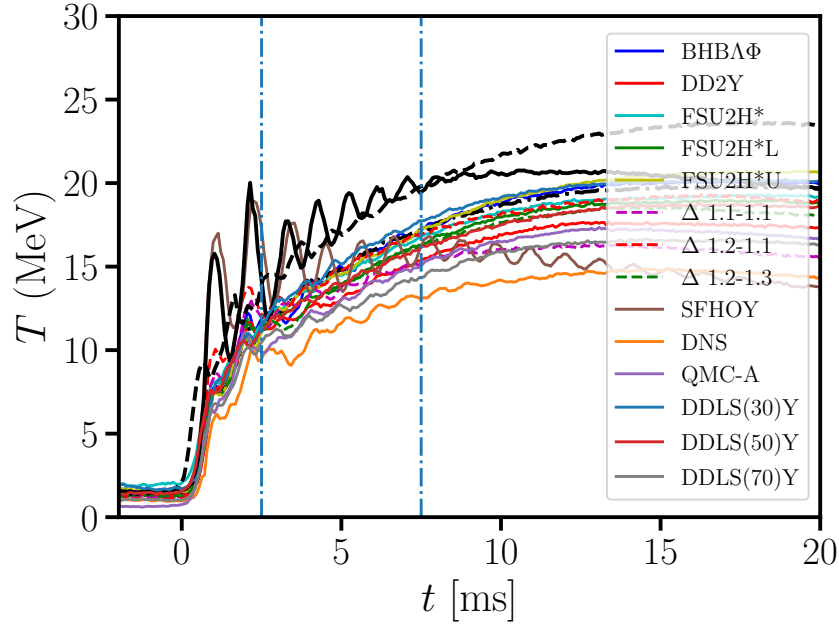


Figure 4.11: Time evolution of the average temperature for different hyperonic EoSs. Three different nucleonic models (solid line for SFHO, dashed line for FSU2R, and dash-dotted line for DD2) are also shown. The figure was published in [369].

Finally, in Fig. 4.11, the mass-averaged temperature evolution of the remnant is shown. Before the merger, the matter is cold, with only minimal numerical heating. After the merger, the temperature rises significantly for all EoSs. For the softest EoSs, strong initial oscillations caused by the dynamic rearrangement and the shock heating are noticeable. In the later phases, the average temperature in the remnant tends to stabilize for all EoSs.

As seen in the figure, the predicted temperatures vary significantly between the models. In general, softer EoSs yield hotter remnants, a trend consistent with prior research (see, e.g., Ref. [392]). This is due to the fact that softer EoSs allow matter to be more compressed, translating more gravitational energy into thermal energy during the event.

It is interesting to note that the appearance of heavy baryons in matter influences the average temperature of the remnant. Indeed, the temperature predicted by the purely nucleonic models is higher than the one predicted by their hyperonic counterparts. The difference is apparent when one analyzes specific pairs of models that share the same underlying nucleonic interaction but differ only by the inclusion of heavy baryons. In Fig. 4.11 such pairs are the models from the DD2Y family and DD2 or SFHOY and SFHO. Despite the fact that the models with hyperons and delta baryons are generally softer (they have smaller radii for a fixed NS mass, once the hyperons have appeared in matter), the average temperatures predicted in simulations employing these models are remarkably lower than those of purely nucleonic models. This seemingly counterintuitive result is directly explained by the significantly larger specific heat at constant pressure for hypernuclear matter, as demonstrated earlier in this chapter (Fig. 4.7). As we have discussed, the presence of additional heavy baryon species provides a greater number of available microstates for energy distribution. The excess thermal energy is absorbed by populating the new hyperonic degrees of freedom rather than solely increasing the kinetic energy of existing nucleons. This acts as an efficient cooling mechanism within the remnant. For the interested reader, we refer to the works of Refs. [392, 393] for related analyses of purely nucleonic models.

3 BNS observables

In this section, we show the results that we have obtained for the merger observables from the simulation campaign. We will again focus on the simulations of $1.4M_{\odot} - 1.4M_{\odot}$ NS mergers. Results for asymmetric and/or heavier BNS mergers will be shown at the end of the section. The most important results are also listed in tabular form in Appendix C.

As earlier mentioned, two simulations were performed with each EoS from both sets (hyperonic and nucleonic). Here we will show the results obtained from the 3D run, and the difference between the results obtained using the 3D EoSs and the constant $\Gamma_{\text{th}} = 1.75$

approach. The results obtained with the latter approach will always have the super-index 1.75, e.g. $X^{1.75}$, for any arbitrary quantity X (see the flowchart in Fig. 4.12).

3.1 Dominant frequency peak f_{peak}

We start by focusing on the results related to the dominant frequency peak f_{peak} , and the so-called frequency shift,

$$\Delta f = f_{\text{peak}} - f_{\text{peak}}^{1.75}, \quad (4.8)$$

with f_{peak} being the dominant GW frequency from the simulation using the 3D EoSs and $f_{\text{peak}}^{1.75}$ being the dominant frequency from the $\Gamma_{\text{th}} = 1.75$ run. We remind the reader that the latter corresponds to a calculation with the zero-temperature beta-equilibrium EoS slices for each model, with a thermal contribution assuming a constant thermal index of $\Gamma_{\text{th}} = 1.75$.

Before we discuss the results, a few words are needed to describe how those peaks are obtained from the simulations. Specifically, we use the strain of the cross polarization, $h_{\times}(t)$, obtained directly from our numerical relativity simulations. The Fourier transform, $h_{\times}(f)$, is then applied to a chosen time window of the post-merger signal, converting the time-domain waveform into its frequency components. We then compute the effective amplitude spectrum, defined as

$$h_{\text{eff},\times}(f) = f \cdot h_{\times}(f). \quad (4.9)$$

This weighting by frequency amplifies the contributions from higher-frequency modes, which are often the most relevant for post-merger dynamics. Since the resulting effective amplitude spectrum $h_{\text{eff},\times}(f)$ consists of discrete data points, a precise determination of the peak's maximum frequency and amplitude requires an interpolation method. With the three highest data points at the top of each peak, a quadratic fit is performed. The position of the maximum of the parabola represents the position of the frequency peak. This method offers a more accurate estimate of the true peak frequency compared to simply selecting the highest discrete data point, which can be sensitive to the frequency resolution of the discrete Fourier transform. A typical output from our simulation campaign is presented in Appendix D.

Relative frequency shift

In Fig. 4.13 we present the results of the relative frequency shift, $\Delta f / f_{\text{peak}}$, as a function of f_{peak} . In the plot, different symbols categorize the different types of EoSs: purely nucleonic models are indicated by plus signs, models incorporating hyperons are marked with crosses, and models that include both hyperons and Δ baryons are represented by asterisks. The color bar provides additional information, indicating the ratio between the

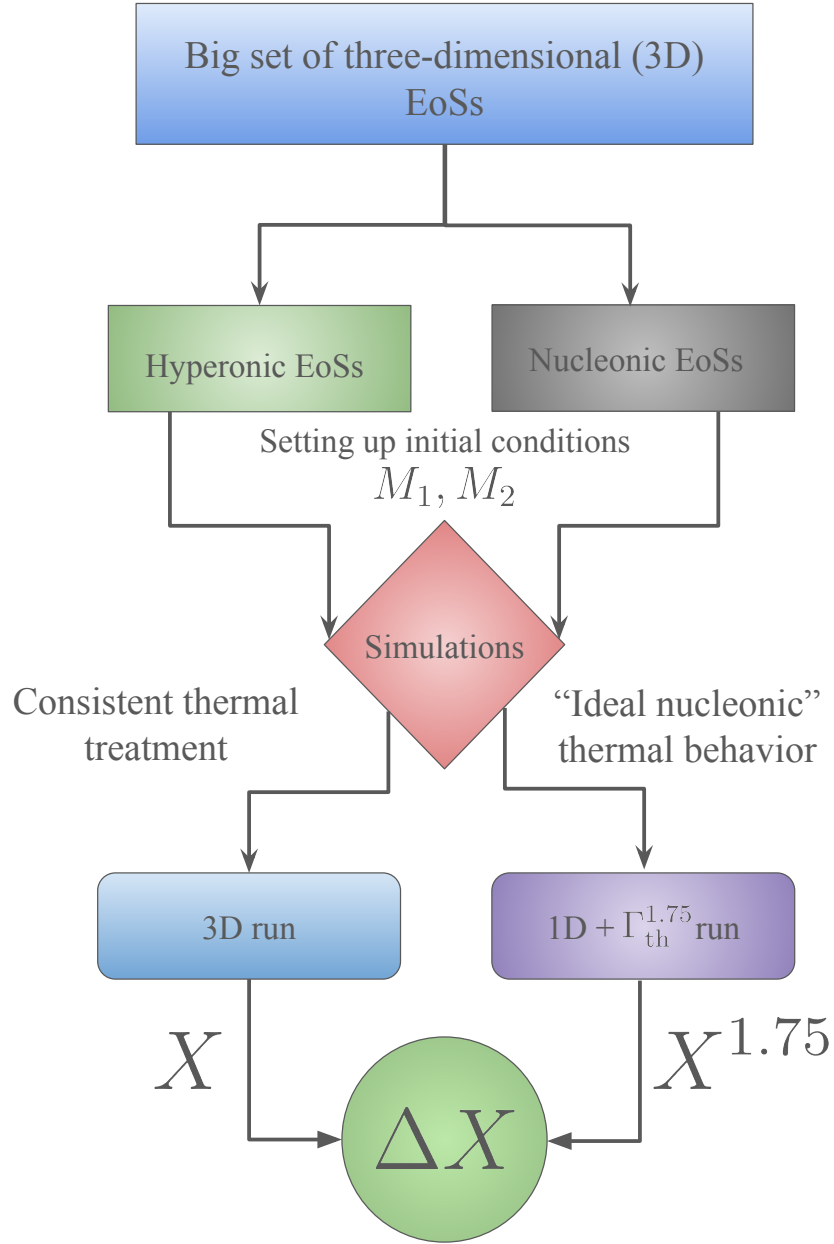


Figure 4.12: Flowchart of the strategy implemented in the simulation campaign. The difference between an observable X obtained from the 3D run and the $\Gamma_{\text{th}} = 1.75$ run will be labeled as ΔX .

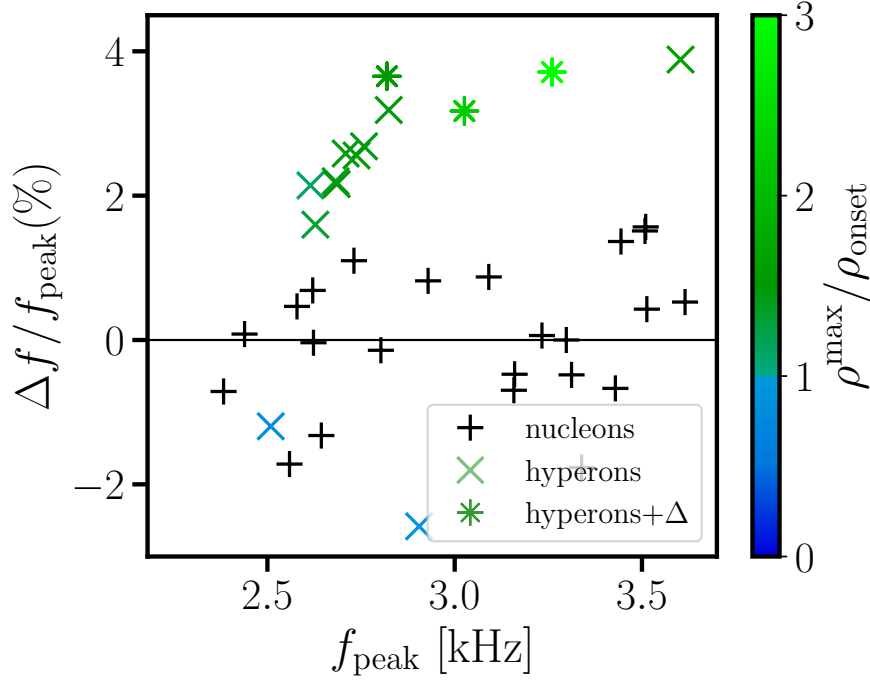


Figure 4.13: Relative frequency shift as a function of the peak frequency for all EoSs from both sets. Purely nucleonic models are indicated by plus signs, models incorporating hyperons are marked with crosses, and models that include both hyperons and Δ baryons are represented by asterisks. The color bar indicates the ratio between the maximum rest-mass density achieved in the merger remnant within the first 5 milliseconds after the merger and the zero-temperature onset density of heavy baryons ($\rho^{\max}/\rho_{\text{onset}}$). The figure was published in [369].

maximum rest-mass density achieved in the merger remnant within the first 5 milliseconds after the merger and the zero-temperature onset density of heavy baryons ($\rho^{\max}/\rho_{\text{onset}}$). This ratio gives an indication of the potential for hyperon population in the remnant.

From Fig. 4.13 it is clear that the purely nucleonic models consistently group around $\Delta f = 0$. Indeed, we confirm that the choice of $\Gamma_{\text{th}} = 1.75$ reproduces well the thermal behavior of nucleonic models. We can see that the quantity Δf can be interpreted as a measure of how much a given model deviates from an “idealized” nucleonic thermal behavior. Conversely, the hyperonic models exhibit a systematic and positive shift towards higher frequencies, typically varying in the range of a few percent. This systematic shift is a direct consequence of the thermal softening of the EoS induced by the appearance of the heavy baryons in the dense, hot post-merger matter. As discussed in the preceding sections, the population of hyperons leads to a reduction in the thermal pressure and a lower effective thermal index, softening the overall EoS. A softer EoS generally

corresponds to more compact merger remnants for a given mass and angular momentum, which, in turn, oscillate at higher frequencies. However, this frequency shift is not universal for all hyperonic models. When the maximum density reached in the remnant is not sufficiently high to induce a significant abundance of heavy baryons, as indicated by a low value on the color bar, the characteristic frequency shift is not observed. In these cases, the hyperonic models yield frequency shifts that are indistinguishable from those of purely nucleonic EoSs. The models that have this behavior, DNS and QMC-A, were already identified in the previous section as models that do not predict a meaningful amount of hyperons to be produced in the remnant with this configuration.

As an additional confirmation that the positive frequency shift that the hyperonic EoSs yield is indeed caused by the thermal softening of the EoS, in Fig. 4.14 the relative frequency shift is shown again, but this time as a function of the mass- and time-averaged thermal index in the remnant, $\bar{\Gamma}_{\text{th}}$, computed for the 3D EoS runs. We calculate $\bar{\Gamma}_{\text{th}}$ by averaging $\Gamma_{\text{th}}^{\text{av}}$ over a time window of 5 ms starting at 2.5 ms after the merger. Indeed,

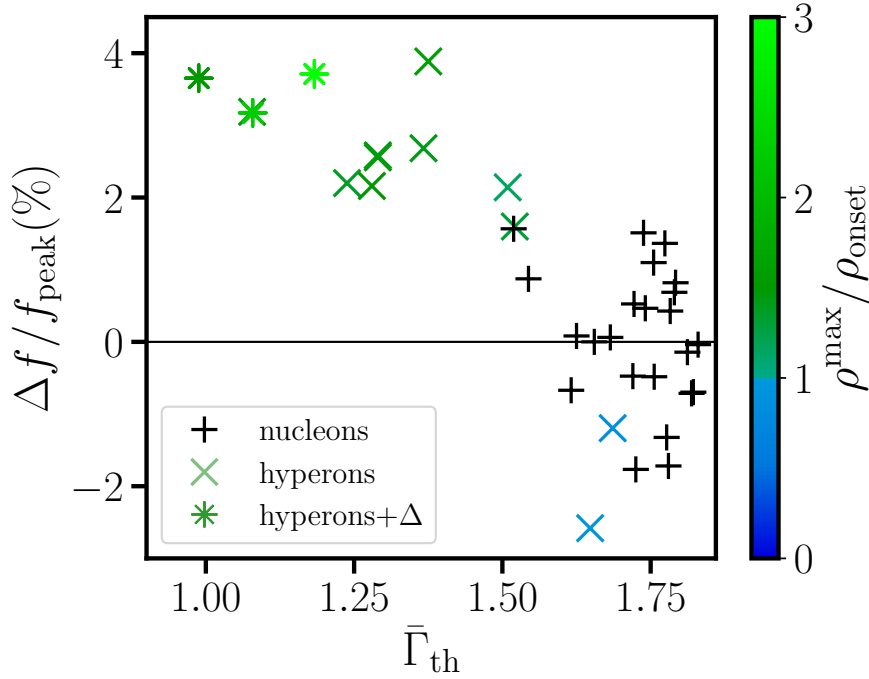


Figure 4.14: Relative frequency shift as a function of the mass- and time-averaged thermal index in the remnant for all EoSs from both sets. The notation is the same as in Fig. 4.13. The figure was published in [369].

our assumed correlation is confirmed: a smaller average thermal index corresponds to a larger frequency shift. Nucleonic models are clustered in the lower-right part of the

figure, while hyperonic models, which have reached the onset density, are mostly found in the upper-left part of the figure.

To get a sense of how much hyperon abundance is needed to be produced in the remnant for a shift to be present, in Fig. 4.15 we again show the relative frequency shift but now as a function of the mass- and time-averaged hyperon fraction in matter $\bar{Y}_{\text{hyp}}$. This quantity is computed analogously to $\bar{\Gamma}_{\text{th}}$, by time-averaging Y_{av} . From the figure it is clear that

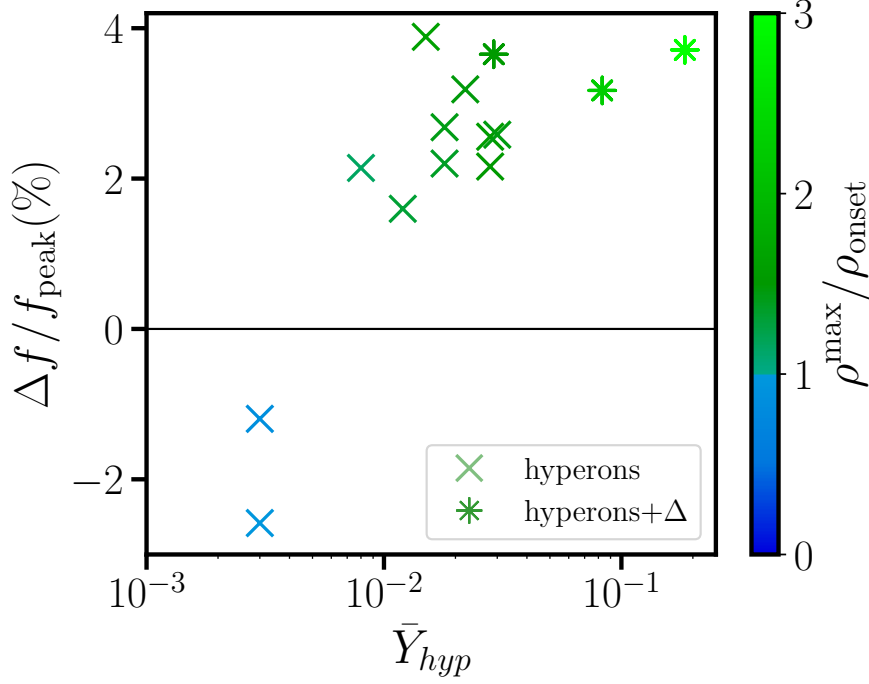


Figure 4.15: Relative frequency shift as a function of the mass- and time-averaged hyperon fraction $\bar{Y}_{\text{hyp}}$ in the remnant. Models that incorporate hyperons are marked with crosses, while models that include both hyperons and Δ baryons are represented by asterisks. The color bar indicates the ratio between the maximum rest-mass density achieved in the merger remnant within the first 5 milliseconds after the merger and the zero-temperature onset density of heavy baryons ($\rho^{\text{max}}/\rho_{\text{onset}}$). The figure was published in [369].

even a relatively small average hyperon abundance of 1%, can induce an increase in the dominant GW frequency. Interestingly, the frequency shift appears to saturate, not exceeding a few percent, even when the total mass-averaged number of hyperons in the remnant continues to increase significantly to much higher fractions. This intriguing phenomenon, where the GW frequency shift shows a plateau despite a growing hyperon population, can be explained by the density-dependent trend of the thermal index and

the relation with the thermally produced hyperons. As we have illustrated (Fig. 4.5), the thermal index at constant temperature and charge fraction in hypernuclear matter exhibits a characteristic recovery at densities far away from the onset density for the particular species. We connected this behavior with the decrease of $\Delta\rho_Y$ (Fig. 4.2). While the overall hyperon population increases with density, the excess of hyperons produced solely due to thermal effects is reduced at very high densities. The softening effect to which our quantity Δf is sensitive is associated with the thermally driven redistribution of degeneracy pressure that diminishes at very high densities. Hence, although the total hyperon fraction may continue to rise, additional thermal softening is not produced in the remnant. This suggests that while hyperons undoubtedly leave an imprint on the GW spectrum, the magnitude of this imprint, specifically the frequency shift, saturates once the EoS has undergone its most significant thermal softening related to the initial appearance of thermally populated exotic degrees of freedom.

Before we finish this part, let us comment on the usefulness of the obtained systematic behavior of the frequency shift and the possibility of detecting it in future observations. The systematic shift of the dominant post-merger GW frequency due to the presence of hyperons ranges from approximately 50 Hz to 150 Hz. Although this shift seems small compared to the intrinsic full width at half maximum of the spectral peaks or to the overall cold EoS stiffness variation of f_{peak} , it is nonetheless potentially sizable enough for a detection by future GW observatories. We envision a scenario where an extracted value f_{peak} from a GW event could serve as a discriminator for the presence of exotic matter. We will suppose that the cold EoS of dense matter could be constrained with high precision through the analysis of many cold star observations, but its composition is still unknown. In such a scenario, one could perform a reference merger simulation using the precisely constrained cold EoS, but assuming a canonical nucleonic thermal behavior, by employing an ideal-gas contribution with a constant thermal index $\Gamma_{\text{th}} = 1.75$. The predicted reference value, extracted from the simulation, $f_{\text{peak}}^{1.75}$, would then serve as a benchmark. A subsequent observed deviation between the predicted $f_{\text{peak}}^{1.75}$ and the measured f_{peak} from a real event, specifically a positive shift of approximately 50 Hz to 150 Hz, would suggest the presence of hyperons in the merger remnant. The feasibility of this inference relies on the observational capabilities of next-generation GW detectors. Recent GW injection and parameter estimation studies (Refs. [394, 395, 396, 397, 398, 399, 400, 401, 402]) indicate that future facilities, such as ET or CE, may be able to determine f_{peak} with a precision of ~ 10 Hz. Therefore, in theory, a frequency change of 50 – 150 Hz is measurable. However, the successful implementation of this strategy faces significant prerequisites. As we stated, the cold EoS must indeed be measured with sufficiently high precision from observations. In addition, the requirement that numerical relativity simulation tools are reliable enough to predict an accurate reference value of f_{peak} for comparison with observational data needs to be fulfilled too. This implies that the uncertainties associated with the numerical

methods, the microphysical input (such as neutrino transport and magnetic fields), and the modeling of the nucleonic EoS itself must be sufficiently small. Both prerequisites are currently not satisfied but may be fulfilled in the following decades.

$f_{\text{peak}} - \Lambda$ relation

We have explained that inferring the presence of heavy baryons in matter through the frequency shift is a difficult task. From an observational standpoint, it is interesting to connect two quantities that can be measured experimentally, which may suggest the actual composition of dense matter. In some recent works, it was shown that the relation between the dominant frequency peak and the tidal deformability of cold neutron stars, the so-called $f_{\text{peak}} - \Lambda$ relation, can be particularly tight for purely nucleonic models, but it can deviate when quark matter is present in the remnant [403, 404]. We will adopt the same strategy and explore if hyperons induce a deviation in this relation. This method of inferring the composition would bypass the requirement to know precisely the cold EoS.

A first step in this analysis is to identify a reference point M where the $f_{\text{peak}} - \Lambda_M$ relation for purely nucleonic models becomes particularly tight. The index M labels the mass of the cold star in β -equilibrium for which the tidal deformability Λ has been used. To achieve this, we systematically analyze for different values of M , the scatter of our nucleonic EoS data points with respect to a quadratic function $f_{\text{fit}}(\Lambda_M)$ [364, 403]. Indeed, for each fixed reference mass M , we quadratically fit the $f_{\text{peak}}(\Lambda_M)$ results and thus obtain the $f_{\text{fit}}(\Lambda_M)$ curves. The quality of this fit is assessed by computing the residual for each data point, defined as $r_i = |f_{\text{peak}} - f_{\text{fit}}|$. We then quantify the overall scatter by determining the average value of these residuals, $r_{\text{avg}} = \sum r_i / N$ (where N is the number of nucleonic data points) and the maximum residual, $r_{\text{max}} = \max(r_i)$.

In Fig. 4.16, we represent the maximum and the average deviation of the data from the quadratic fit as functions of the reference stellar mass M . From the figure it is clear that while the average scatter with respect to a quadratic fit, r_{avg} , is hardly sensitive to M , the maximum deviation, r_{max} , exhibits a minimum at a reference mass of $M = 1.65 M_{\odot}$. This minimum indicates that at this particular mass, the quasi-universal relation for nucleonic matter shows the tightest correlation, minimizing the model-dependent scatter. Consequently, we chose $M = 1.65 M_{\odot}$ as our reference mass for the subsequent comparison with the hyperonic models.

We provide the $f_{\text{peak}} - \Lambda_{1.65}$ relation for all EoSs taken into consideration in this work in Fig. 4.17. The notation for the symbols remains consistent with the previous figures in this chapter. In addition, the solid black line represents the quadratic fit obtained from the purely nucleonic EoS data, serving as our nucleonic baseline, while the maximum residual of these nucleonic models with respect to the fit is labeled with the surrounding

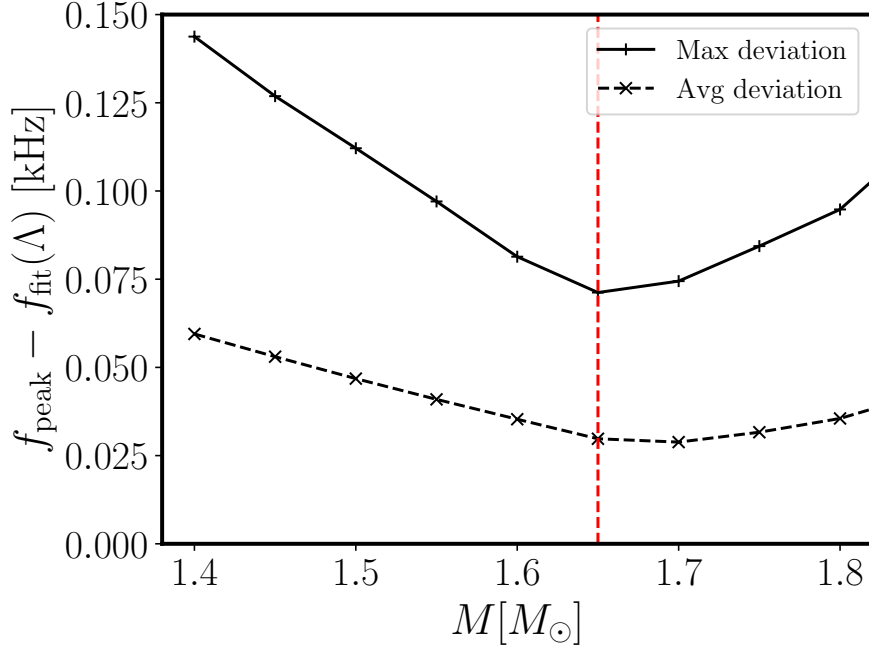


Figure 4.16: Average and maximum deviation for purely nucleonic EoSs from quadratic least-squares fits $f_{\text{peak}} - \Lambda_M$. The solid line shows the behavior of the maximum deviation as a function of the reference mass, while the dashed line shows the functional dependence of the average deviation. The point M at which the maximum deviation reaches a minimum is labeled with a red vertical line. The figure was published in [369].

gray band. This band thus defines the expected "nucleonic corridor" within which a purely nucleonic system is anticipated to lie, accounting for the present model uncertainties.

It is noticeable that for a fixed $\Lambda_{1.65}$, f_{peak} of most models that incorporate hyperons and Δ resonances are systematically higher compared to that predicted by the nucleonic EoSs. The ones that do not follow this trend are those that do not produce a significant amount of hyperons in the remnant, as we have discussed before. We reiterate that the upward shift is a direct consequence of the EoS softening induced by the appearance of hyperons and/or Δ baryons in the dense and hot merger remnant. Some models featuring heavy baryons lie distinctly outside the defined gray nucleonic corridor. This observable separation provides a means to distinguish hypernuclear matter from purely nucleonic compositions. An observed $f_{\text{peak}} - \Lambda_{1.65}$ data point that falls outside the nucleonic band, particularly towards higher frequencies, would suggest that the heavy baryons are present in the post-merger phase of a BNS merger.

A few remarks have to be addressed too. The effects that the appearance of hyperons have on the dominant post-merger frequency are similar to the ones predicted by the possible

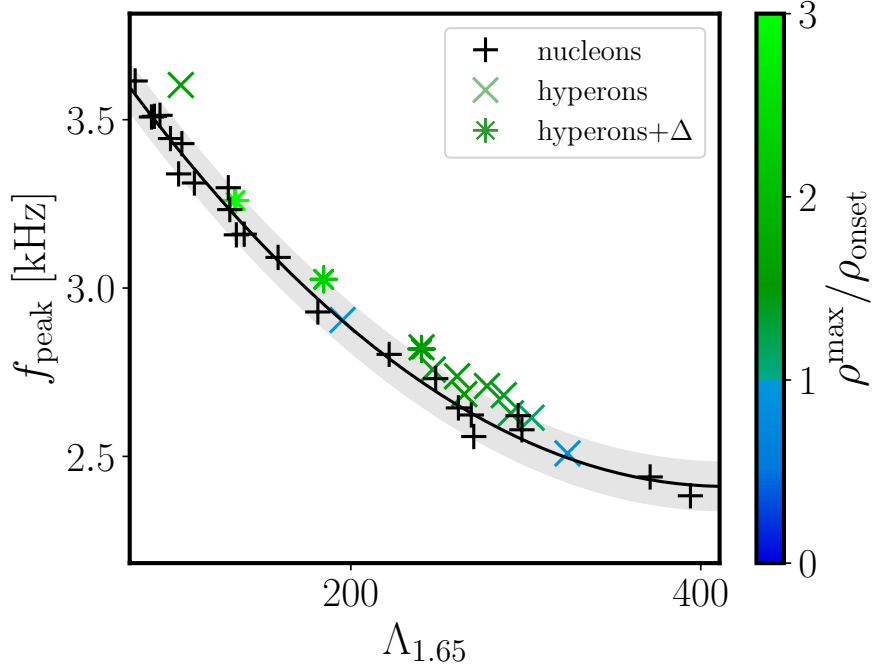


Figure 4.17: $f_{\text{peak}} - \Lambda_{1.65}$ relation for all EoSs considered in this work. The notation of the symbols and the color bar is the same as in Fig. 4.13. The solid black line represents the quadratic fit obtained from the purely nucleonic EoS data. The gray band represents the maximum residual of these nucleonic models with respect to the fit. The figure was published in [369].

appearance of quark matter [405, 404]. However, other observables like the mass-radius of the stars might be sensitive to the appearance of quark matter. Furthermore, some nucleonic models predict that the thermal behavior of purely nucleonic matter can mimic the behavior of the hypernuclear matter [406, 407, 408]. We specifically addressed this issue in Appendix E.

$f_{\text{peak}} - \rho_{\text{max}}$ relation

The relationship between f_{peak} and the maximum rest-mass density achieved in the merger remnant within the first 5 milliseconds after the merger, ρ_{max} , is also intriguing. It can be particularly important as it offers a direct pathway to infer the extreme densities reached in the remnants from an observable standpoint. The relation has been studied in some previous works [404, 405], where it has been concluded that it might be sensitive to the nature of phase transitions within the remnant, specifically exhibiting deviations from the otherwise almost-linear trend that characterizes nuclear matter for EoSs that feature a strong first-order phase transition to quark matter.

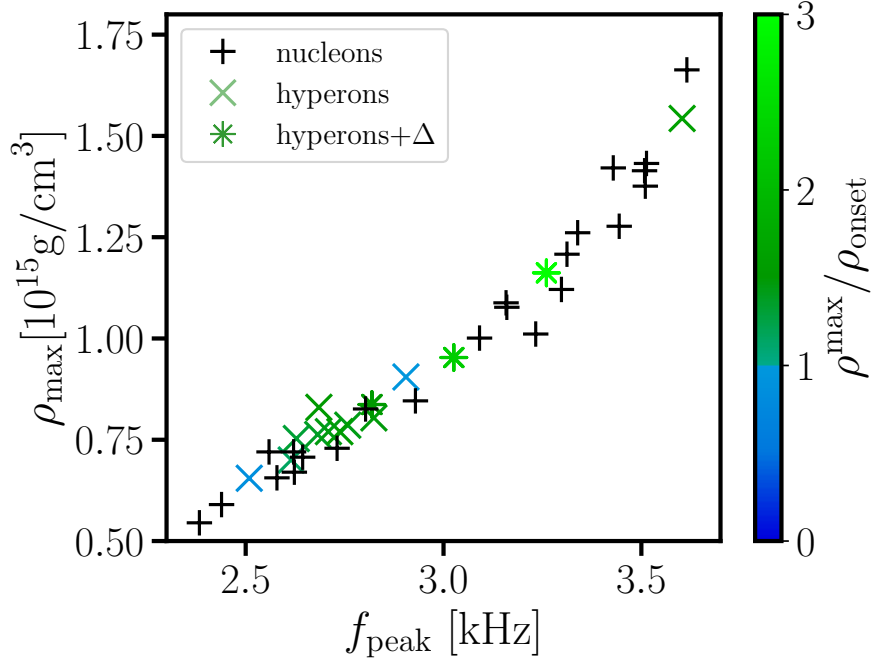


Figure 4.18: $f_{\text{peak}} - \rho_{\text{max}}$ relation for all EoSs considered in this work. The notation is the same as in Fig. 4.13. The figure was published in [369].

In Fig. 4.18, we show the $f_{\text{peak}} - \rho_{\text{max}}$ points for all models employed in the thesis. The analysis reveals that there are no significant differences between the relation for the purely nucleonic models and those including exotic baryonic degrees of freedom. Unlike strong first-order phase transitions, the gradual onset and population of heavy baryons do not fundamentally break this quasi-universal correlation. Therefore, we conclude that this robust correlation holds for the entire class of hadronic models.

The universality of this relation implies that it can serve as a tool to estimate the maximum rest-mass density reached in the remnant solely from an observation of f_{peak} , even if the presence of heavy baryons is initially unknown. It is possible to use this fact to constrain the onset density of heavy baryons.

Specifically, if a future GW event indicates that there are no hyperons in the remnant (e.g., if the observed $f_{\text{peak}} - \Lambda_M$ point lies firmly within the nucleonic corridor, as discussed previously), then, by using the universal $f_{\text{peak}} - \rho_{\text{max}}$ relation, the maximum density attained in that particular merger can be inferred. This density would then effectively place a lower bound on the density at which hyperons are still not present in matter (in significant, detectable amounts).

Conversely, if a GW event indicates that hyperons/ Δ -baryons are present in the remnant (i.e., the $f_{\text{peak}} - \Lambda_M$ point falls outside the nucleonic band), the same universal relation would again allow us to determine the maximum density reached. The inferred density would then effectively serve as an upper bound for the onset density at which heavy baryons are produced in matter in significant, detectable amounts.

Direct comparison of EoSs with and without hyperons

In the first works investigating the impact of hyperons in BNS merger events, the strategy revolved around using one chosen EoS that predicts the appearance of hyperons in the dense matter (Refs. [409, 410]). Two simulations were then performed: one with the chosen hyperonic EoS and one with an EoS with the same nucleonic interaction as the hyperonic one, which artificially prohibited the appearance of hyperons in matter. The differences in the observables obtained from the two simulations were then attributed to the appearance of hyperons.

For completeness, in this subsection we also investigate the impact of hyperons for a given nucleonic interaction. We identify specific subsets of our EoS models, which consist of hyperonic models that have a purely nucleonic counterpart. This comparison allows us to disentangle the effects arising from the exotic degrees of freedom from those inherent to the nucleonic EoS. Within this subset, a few models are built within the nucleonic DD2 interaction: DD2Y, BHBA ϕ , DD2Y Δ 1.1 – 1.1, DD2Y Δ 1.2 – 1.1, and DD2Y Δ 1.2 – 1.3. These hyperonic EoSs differ in the specific heavy species included and in the predicted YY and YN interactions. The SFHO nucleonic EoS has only one hyperonic counterpart, the SFHOY model. Finally, each of the nucleonic models from the DDLSN family has a corresponding DDLSY hyperonic equivalent

In Fig. 4.19, the dominant post-merger GW frequency is presented as a function of the radius (top panel) and the tidal deformability (bottom panel) for the selected subset of EoS models and for two representative NS masses, $1.4 M_{\odot}$ (left panel) and $1.8 M_{\odot}$ (right panel). Lines are used to explicitly connect models that are based on the identical nucleonic interaction, providing a visual representation of the quantitative impact of the inclusion of hyperons on f_{peak} and the stellar properties. Unlike in the other plots of this chapter, the color bar in the figure illustrates the time-averaged fraction of heavy baryons inside the remnant, computed in a 2.5 ms window 5 ms after the merger.

The inclusion of hyperons and/or other heavy baryons consistently leads to an increase of the dominant frequency peak. This upward shift can be substantial, amounting to more than 0.5 kHz in some models, significantly exceeding the shifts discussed in the context of general nucleonic vs. hyperonic comparisons. At the same time, the radius and tidal deformability for a given mass are also reduced with the inclusion of heavy baryons, reflecting the softening of the cold EoS. This reduction in radius can be greater than 1 km

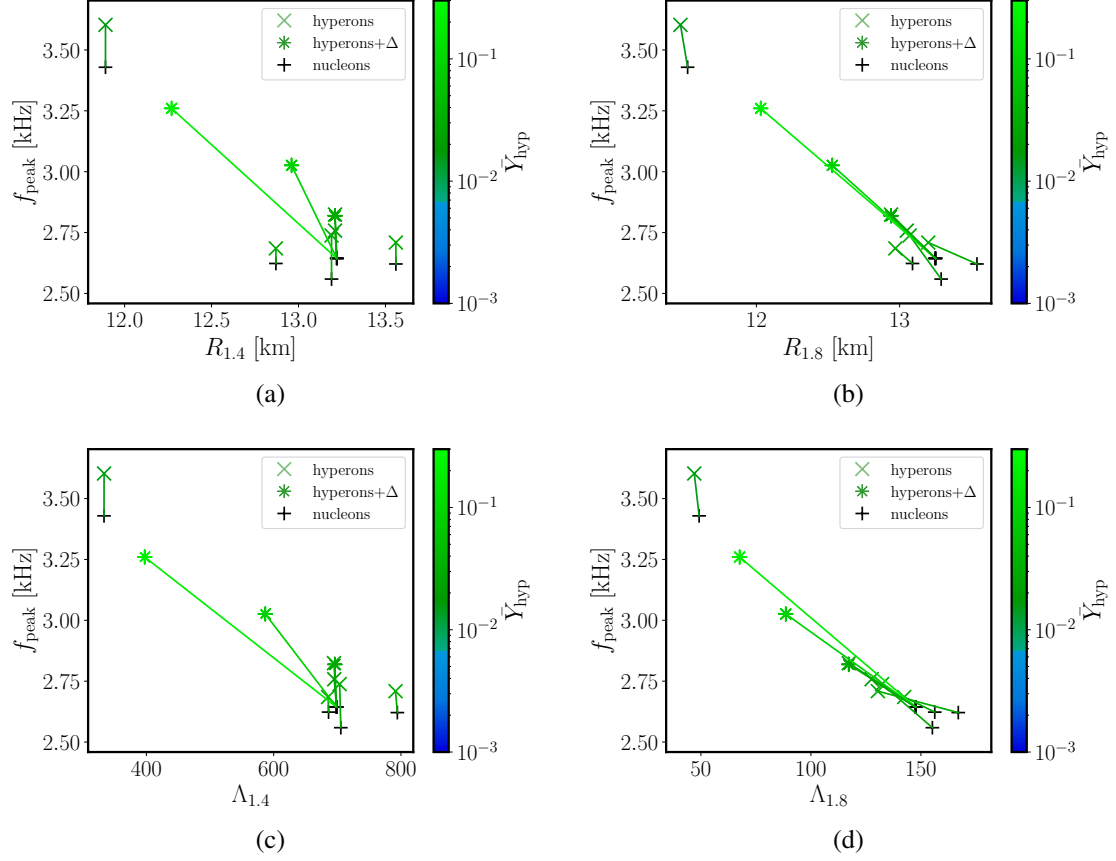


Figure 4.19: f_{peak} as a function of stellar parameters of cold nonrotating NS: a) $R_{1.4}$, b) $R_{1.8}$, c) $\Lambda_{1.4}$ and d) $\Lambda_{1.8}$, for the chosen subset of nucleonic and hyperonic models that share the same nucleonic interaction. Purely nucleonic models are indicated by plus signs, models incorporating hyperons are marked with crosses, and models that include both hyperons and Δ baryons are represented by asterisks. The color bar shows the fraction of hyperons in the remnant. The figure was published in [369].

for certain stellar masses (panel b), and the tidal deformability can be reduced by a factor of 2 (panel d), indicating a more compact stellar structure. However, it is noteworthy that for the properties relevant to lower-mass stars, which are presented in the panels a) and c), there may not be an effect from the hyperons. Hyperons/ Δ -baryons, for most EoSs, typically appear in substantial quantities only at densities exceeding the central density of a $1.4M_{\odot}$ mass NS. Therefore, their presence does not significantly alter the properties of such lower-mass, cold stars.

We must emphasize that this straightforward, side-by-side comparison between nucleonic and hyperonic models sharing identical nucleonic interactions is primarily of academic value, as this comparison is likely not directly useful for deducing the presence of hyperons from actual observed BNS merger events. As previously explained, a perfectly corresponding purely nucleonic model to compare against a real observation does not exist in nature, given that the 'true' EoS is unique. However, one might hope for a future scenario where highly precise theoretical model, potentially backed up by robust experimental data from terrestrial nuclear physics facilities, could accurately describe a hypothetical system composed only of nucleons. When compared to an actual observation of f_{peak} or inferred stellar parameters from GW and electromagnetic observations, a measurable discrepancy could be found. If this discrepancy is consistent with the trends presented here, it would provide strong evidence that heavy baryons exist in the remnant matter.

Numerical robustness

We have seen that the GW shifts induced in matter by thermally populated hyperons are of the order of 50-150 Hz. This raises the question of whether the phenomenon could be a numerical artifact of the hydrodynamic scheme used. To explicitly demonstrate that the shift is indeed a physical feature, we conducted additional simulations with another numerical scheme, the moving-mesh AREPO code [411, 412, 413]. A main property of this scheme is that the numerical artifacts across grid cells, as a consequence of the advecting matter, are minimized. The AREPO code, like our original solver, also utilizes the CFA to solve the Einstein equations.

The simulations with the AREPO code used DD2 and DD2Y EoSs, both in a full 3D and in a $\Gamma_{\text{th}} = 1.75$ run. We report the results in Table 4.1, which are compared to the results obtained with the SPH scheme. One can easily see that the values for the frequency peak for the DD2 EoS show good agreement between both simulations. Interestingly, when the DD2Y EoS is used, a small discrepancy can be observed. Nevertheless, we note that the shift Δf_{peak} consistently appears in the AREPO simulations as well and is of roughly the same magnitude as the one that we have obtained with the SPH simulations. Indeed, a positive shift of 70 Hz is predicted by the AREPO simulation with the DD2Y EoS. This

confirms that the frequency shift we have found is not a numerical artifact but is indeed a physical feature induced by the thermal hyperons.

	DD2		DD2Y	
	$f_{\text{peak}}[\text{kHz}]$	$f_{\text{peak}}^{1.75}[\text{kHz}]$	$f_{\text{peak}}[\text{kHz}]$	$f_{\text{peak}}^{1.75}[\text{kHz}]$
AREPO	2.63	2.67	2.71	2.64
SPH	2.64	2.68	2.82	2.73

Table 4.1: AREPO and SPH results of the dominant GW frequency peak using the DD2 EoS and the DD2Y EoS. The table was published in [369].

To ensure the robustness of our findings and assess the numerical convergence of our simulations, we also briefly investigated the influence of spatial resolution on the extracted GW frequencies. The analysis was conducted by performing simulations with the SPH code for several different EoSs employing three distinct particle resolutions: approximately 150,000, 300,000, and 600,000 SPH particles, respectively. This allows us to gauge the sensitivity of our results to numerical discretizations. The results are shown in Table 4.2.

EoS	$6 \cdot 10^5$ particles		$3 \cdot 10^5$ particles		$1.5 \cdot 10^5$ particles	
	$f_{\text{peak}}[\text{kHz}]$	$f_{\text{peak}}^{1.75}[\text{kHz}]$	$f_{\text{peak}}[\text{kHz}]$	$f_{\text{peak}}^{1.75}[\text{kHz}]$	$f_{\text{peak}}[\text{kHz}]$	$f_{\text{peak}}^{1.75}[\text{kHz}]$
DD2	2.623	2.678	2.644	2.679	2.635	2.692
SFHO	3.430	3.396	3.429	3.452	3.435	3.495
DD2Y	2.789	2.697	2.824	2.734	2.835	2.764
DD2Y $\Delta 1.2-1.3$	3.230	3.121	3.260	3.139	3.297	3.205
SFHOY	3.538	3.485	3.603	3.463	3.615	3.537

Table 4.2: Summary of the dominant frequency peak values obtained for the numerical convergence study.

Across this set of resolutions, the absolute values of the dominant post-merger GW frequencies, f_{peak} , change by at most a few tens of Hz. More specifically, for the nucleonic runs, we observe stable results with increasing resolution, especially when the 3D EoSs are used. For example, the 3D run for the SFHO model yielded a difference of less than 10 Hz between the different resolutions. Similar qualitative results are also obtained with the DD2 EoS. This high degree of stability for nucleonic models suggests good numerical convergence and a lesser sensitivity to the subtle thermodynamic conditions within the remnant.

In contrast, hyperonic models demonstrate a slightly larger, though still controlled, resolution dependence, with a general trend to yield higher absolute frequencies for smaller SPH particle numbers. This difference typically amounts to approximately a few tens of Hz between the runs with the highest and lowest resolution for the models DD2Y, SFHOY, and DD2Y Δ 1.2-1.1. This enhanced sensitivity in hyperonic models can be physically attributed to the more complicated behavior of their EoS.

Despite these minor variations in absolute frequencies across resolutions, we highlight that the relative shifts in f_{peak} due to the presence of hyperons remain consistent. When comparing simulations between runs with different thermal treatments, the characteristic shift of the order of 50-150 Hz due to the presence of hyperons persists consistently throughout all tested resolutions for the hyperonic models. This demonstrates that the influence of hyperons on the post-merger GW signal, discussed in this study, is a reasonably strong physical characteristic.

3.2 Secondary frequency peaks

Beyond the well-studied dominant frequency peak in the GW spectra from BNS merger events, other, often referred to as secondary spectral features in the kilohertz range have been identified in numerical simulations. Although typically weaker in amplitude, these features could carry complementary information about the post-merger dynamics and the EoS of matter. Their detection could become a reality with the future generation of GW observatories (e.g., [313, 414, 415, 416, 417]).

Among these, secondary peaks at frequencies below the dominant frequency are often considered to be more important for potential detection because of instrumental sensitivity. We will focus on two distinct features: the f_{2-0} secondary peak and the f_{spiral} secondary peak. The f_{2-0} spectral feature is typically associated with the lowest-frequency quadrupole oscillation of the remnant, often characterized as a fundamental f -mode or a quasi-radial oscillation, but distinct from the main f_{peak} . The f_{spiral} spectral feature, on the other hand, is believed to originate from the spiral motion of the tidal bulges that form on the surface of the remnant due to differential rotation. For detailed nomenclature and further discussion, see Refs. [391, 418].

Our primary objective in this subsection is to inspect whether thermally produced hyperons leave a discernible signature on these specific secondary GW observables. To investigate this, we adopt an analogous approach to our analysis of f_{peak} . For every EoS in our simulation campaign, we quantify the influence of the different thermal behaviors by defining relative frequency shifts for the secondary peaks:

$$\frac{\Delta f_{2-0}}{f_{2-0}} = \frac{f_{2-0} - f_{2-0}^{1.75}}{f_{2-0}}, \quad \frac{\Delta f_{\text{spiral}}}{f_{\text{spiral}}} = \frac{f_{\text{spiral}} - f_{\text{spiral}}^{1.75}}{f_{\text{spiral}}}. \quad (4.10)$$

In these expressions, we follow the nomenclature we have used in the previous subsection; f_X denotes the frequency of the spectral feature X obtained from simulations employing the full 3D EoS tables. Conversely, $f_X^{1.75}$ represents the frequency obtained from the run with an approximate thermal treatment that mimics an idealized ‘nucleonic’ thermal behavior employing $\Gamma_{\text{th}} = 1.75$. The above ratios, thus, quantitatively capture the deviation of the secondary frequency peaks that directly result from the produced thermal hyperons in the remnants.

The influence of the different thermal behavior on f_{2-0} is presented in the left panel of Fig. 4.20, which illustrates the relation between the relative frequency shift, $\Delta f_{2-0}/f_{2-0}$ and f_{peak} . Analogously to the effect on the dominant peak, hyperonic models that produce a sizable amount of exotic baryons exhibit systematic positive frequency shifts of the f_{2-0} frequency. The relative size of the shift can be even larger than the one observed in the dominant peak, reaching almost 8% for some EoSs. In contrast, purely nucleonic models are clustered around a near-zero shift in $\Delta f_{2-0}/f_{2-0}$. However, they do show an overall larger variation compared to what was observed for the dominant frequency peak. This increased scatter for secondary modes is a common characteristic, often attributed to their generally weaker amplitudes, more complex excitation mechanisms, and potentially higher sensitivity to numerical noise or subtle model differences that might be averaged out in the dominant mode.

The right panel of Fig. 4.20 demonstrates that, as for the dominant peak, we observe an anti-correlation between the shift of f_{2-0} and the average thermal index. Models with a lower $\bar{\Gamma}_{\text{th}}$ tend to exhibit larger positive shifts in f_{2-0} . In contrast to the clear imprints

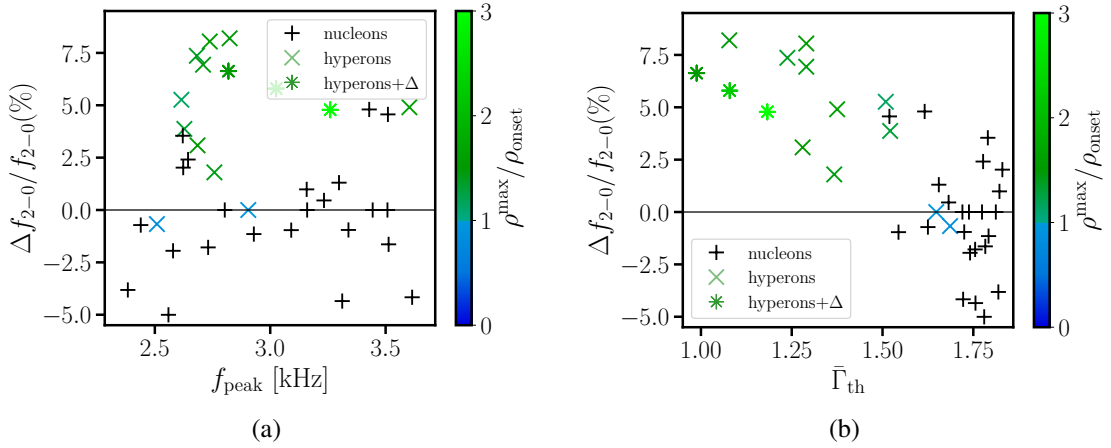


Figure 4.20: $\Delta f_{2-0}/f_{2-0}$ as a function of f_{peak} (left panel) and $\bar{\Gamma}_{\text{th}}$ (right panel). The notation is the same as in Fig. 4.13. The figure was published in [369].

observed on f_{peak} and f_{2-0} , we find a distinct behavior for f_{spiral} , represented in Fig. 4.21.

Thermally produced hyperons and Δ -baryons do not appear to have a significant effect on this particular spectral feature. As shown in Fig. 4.21, the relative frequency shifts for hyperonic models, $\Delta f_{\text{spiral}}/f_{\text{spiral}}$, are largely comparable to those observed for purely nucleonic models and are not affected by the value of f_{peak} (left panel) or $\bar{\Gamma}_{\text{th}}$ (right panel). This indicates that the presence of thermally produced baryons does not induce a systematic or significant deviation in f_{spiral} .

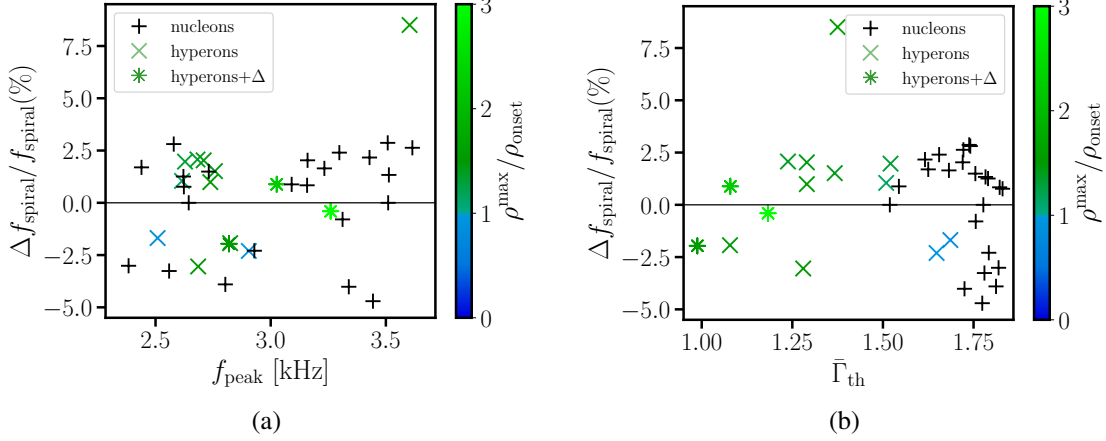


Figure 4.21: $\Delta f_{\text{spiral}}/f_{\text{spiral}}$ as a function of f_{peak} (left panel) and $\bar{\Gamma}_{\text{th}}$ (right panel). The notation is the same as in Fig. 4.13. The figure was published in [369].

The qualitatively different behavior exhibited by f_{peak} and f_{2-0} in contrast to f_{spiral} can be understood by the distinct physical mechanisms responsible for generating these peaks in the first place. As discussed previously, f_{peak} is primarily associated with the fundamental quadrupolar oscillation mode, while f_{2-0} is generated by the quasi-radial oscillation and the quadrupolar mode of the remnant, both of which are sensitive to the finite-temperature EoS. In contrast, f_{spiral} is produced by the orbital motion of two distinct tidal bulges. These bulges originate from the first phases of the merger, just after coalescence. Being predominantly a surface phenomenon and produced in matter that is not significantly heated, f_{spiral} is less sensitive to the later thermal effects induced by hyperon production in the dense matter of the remnant.

To finish this subsection, we also present how the frequency values of the secondary peaks correlate with the value of the dominant frequency peak (Fig. 4.22). We see that indeed there is a strong linear correlation for both secondary peaks. The correlation is not sensitive to the exotic baryons present in matter, as both nucleonic and hyperonic EoSs seem to follow the relation closely.

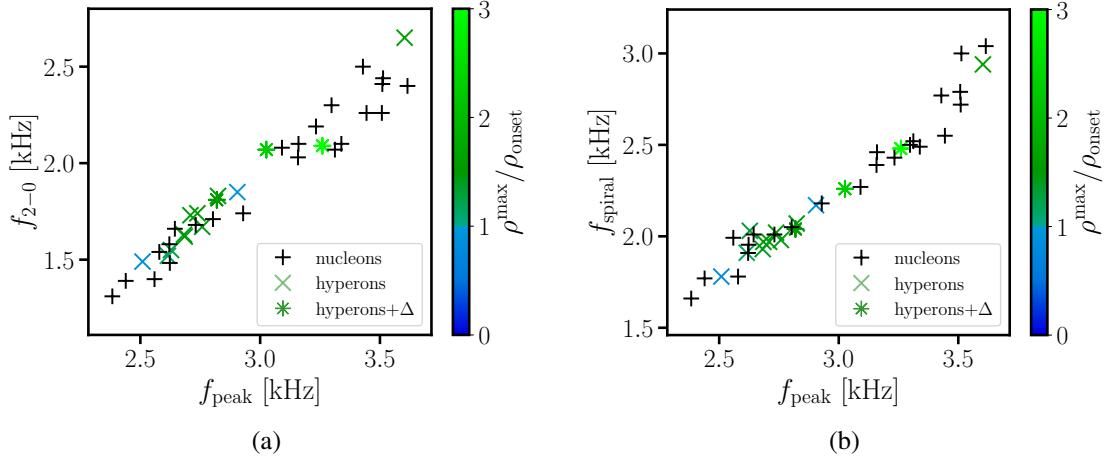


Figure 4.22: Functional dependence between f_{2-0} (left panel) and f_{spiral} (right panel) with f_{peak} . The figure has been published in [369]

3.3 Mass ejecta

Our focus now shifts to another BNS merger observable, the mass ejection, specifically the potential differences in the amount of unbound matter between hyperonic and purely nucleonic EoSs.

We have mentioned that the mass ejection is a particularly important quantity due to its role in the r-process. This process, which occurs in the expanding and cooling phases of the ejecta, is believed to be the dominant cosmic site that produces roughly half of the elements heavier than iron, including all rare-earth elements [52, 419, 420, 421, 422, 423, 424, 425, 426, 427, 428, 429]. The radioactive decays of these newly synthesized unstable isotopes subsequently power a bright electromagnetic transient known as a kilonova (or macronova), which can be observed in the optical and infrared spectrum. The total ejecta mass directly determines the overall amount of synthesized elements and dictates the brightness and timescale of the associated electromagnetic emission.

In our simulations, the amount of ejecta is determined at 10 ms after the merger, a time frame chosen to capture the immediate, so-called dynamical mass ejecta. These ejecta are produced by tidal torques and shock-driven processes [430]. It is important to note that a significant amount of additional ejecta can be produced on much longer timescales, extending up to seconds, through mechanisms such as viscous or magnetically driven outflows from the post-merger accretion disk [420]. These later-time processes are beyond the scope of our current, short numerical simulations. However, the dynamical ejecta are expected to be most directly affected by the properties of the underlying EoS, including the presence of heavy baryons.

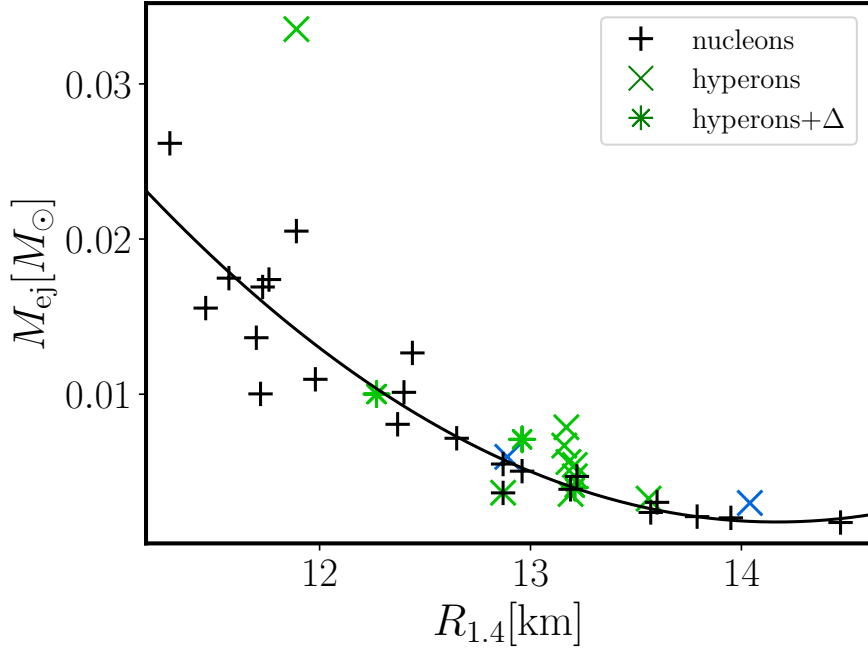


Figure 4.23: Dynamical mass ejecta M_{ej} as a function of the radius $R_{1.4}$ of $1.4 M_{\odot}$ cold NS, 10 ms after merging. Purely nucleonic models are indicated by plus signs, models incorporating hyperons are marked with crosses, and models that include both hyperons and Δ baryons are represented by asterisks. The coloring of the data points signals if the heavy baryons represent more than 1% of the baryon population (green points) or not (blue points). Nucleonic models are labeled with black pluses. The black solid line represents a second-order polynomial fit to the purely nucleonic models. The figure was published in [369].

Fig. 4.23 presents the dynamical ejecta mass obtained from our simulations of equal-mass $1.4\text{--}1.4 M_{\odot}$ binaries as a function of the radius $R_{1.4}$ of the $1.4 M_{\odot}$ cold NS. The radius $R_{1.4}$ serves as a rough measure of the overall stiffness of the EoS at the relevant densities. Smaller radii correspond to softer EoSs, and vice versa. We stick with the notation for the data points as used in previous plots: black plus signs denote purely nucleonic EoSs, while the colored crosses/asterisks represent models with heavy baryons. Specifically, blue symbols indicate models where the time-averaged number fraction of heavy baryons inside the remnant is less than one percent, and green data points highlight models with more than one percent heavy baryons. From the plot it is clear that stiffer EoSs, characterized by larger $R_{1.4}$, tend to produce significantly smaller ejecta masses upon merger [431, 432]. This behavior is primarily attributed to the fact that stiffer EoSs lead to less compact NSs and less compact merger remnants. These remnants do not

undergo vigorous oscillations, which would otherwise facilitate tidal and shock-driven mass ejection.

The majority of hyperonic EoSs generally obey this established trend of purely nucleonic simulations. However, a few of the hyperonic EoSs predict ejecta masses that are substantially increased, up to almost 150%, with respect to nucleonic EoSs with similar stiffness. This is much better visible in Fig. 4.24, where we show the relative difference between the mass ejecta of all models and the one obtained from the quadratic nucleonic fit represented by the solid black line in Fig. 4.23.

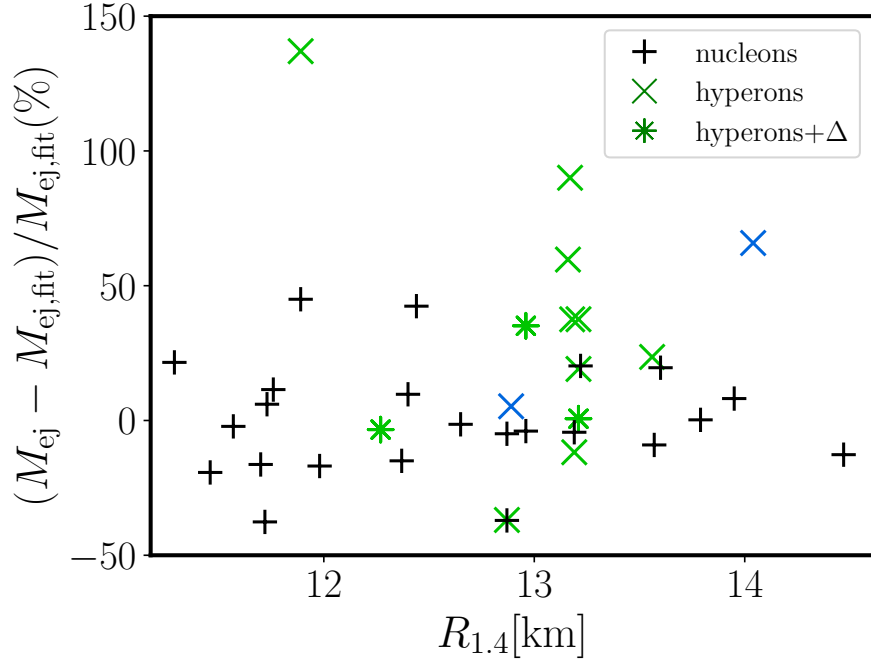


Figure 4.24: Relative difference between M_{ej} and a parabolic fit $M_{\text{ej}}^{\text{fit}}(R_{1.4})$ to the purely nucleonic models. The figure was published in [309].

Even more intriguing is the comparison between the mass ejecta obtained using the full 3D EoSs and the EoSs supplemented with the $\Gamma_{\text{th}} = 1.75$ approximate treatment. In an analogous way as for the frequency peaks, we can investigate the relative difference:

$$\frac{\Delta M_{\text{ej}}}{M_{\text{ej}}} = \frac{M_{\text{ej}} - M_{\text{ej}}^{1.75}}{M_{\text{ej}}}. \quad (4.11)$$

The relative difference between the mass ejecta as a function of $R_{1.4}$ is shown in Fig. 4.25. We can see that the thermal treatment can substantially influence the predicted mass ejecta. However, it is unclear whether a softer/stiffer thermal treatment yields a lower/higher

amount of mass ejecta. Furthermore, the exact composition of the matter does not seem to have a particular influence on the relative difference either.

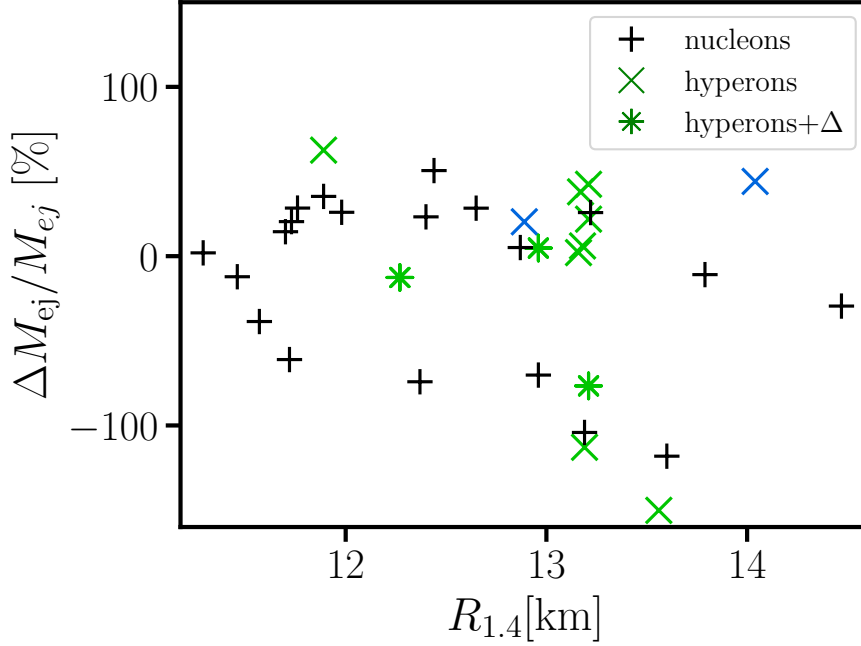


Figure 4.25: Relative difference between the mass ejecta obtained in the 3D run and the one obtained in the $\Gamma_{th} = 1.75$ run. The notation is the same as in Fig. 4.23.

We conclude that translating the theoretical prediction for the mass ejecta into a definitive observational inference is still not possible. A more comprehensive understanding of mass ejection and kilonova emission in general is necessary, as these phenomena are subject to various theoretical and observational uncertainties.

3.4 Threshold mass

The ultimate fate of the post-merger remnant in a BNS merger depends on the initial total mass of the system. As this mass increases, the lifetime of the remnant, defined as the time until its eventual collapse to a black hole, gradually decreases⁴. If the initial mass is high enough, it is possible for the remnant to undergo a prompt black hole collapse.

The minimum initial mass of the system at which the remnant collapses immediately after merging is known as the threshold mass M_{thres} . It is strongly dependent on the underlying

⁴Note that if the initial masses of the stars are low enough, the object can eventually evolve into a massive NS.

EoS. In a simulation campaign, it is determined by considering a set of simulations with varying total binary masses. We again restrict ourselves to symmetric binary simulations. We first identify M_{stable} as the highest total mass for which the remnant does not collapse promptly to a black hole (i.e., it survives for a timescale significantly longer than about a millisecond after the merger). We then obtain M_{unstable} , which represents the lowest total mass for which the remnant directly forms a black hole within a few milliseconds of the merger. The threshold binary mass is then defined as the average value of these quantities:

$$M_{\text{thres}} = \frac{M_{\text{stable}} + M_{\text{unstable}}}{2}. \quad (4.12)$$

In the past, the EoS dependence of the threshold mass has been studied in detail. It has been shown that it correlates well with some cold stellar properties that characterize the EoS. In particular, it was discussed that observations of the threshold mass can put tight constraints on some cold NS properties, such as Λ , R and M_{max} [433, 434, 435, 436, 437, 438, 439, 440, 441]. The idea has been already used to constrain M_{thres} for the progenitor system of the GW170817 event and thus to constrain the EoS of dense matter [434, 442].

However, the aforementioned works did not include a significant amount of hyperonic EoSs, nor did they explore in detail whether the relations found deviate from those of the heavy baryons. For example, in Ref. [437], a small sample of hyperonic EoSs was used, and the analysis did not show any deviation in the bilinear relations such as:

$$M_{\text{thres}}(M_{\text{max}}, R_{1.6}) = aM_{\text{max}} + bR_{1.6} + c. \quad (4.13)$$

In this subsection, we will revisit whether hyperons alter the universal relation given in Eq. (4.13). Furthermore, we will inspect how different thermal behavior affect M_{thres} . With each EoS at our disposal, we calculate M_{thres} for the two runs that we have been using extensively throughout this chapter: the 3D run and the $\Gamma_{\text{th}} = 1.75$ run.

In Fig. 4.26, we present M_{thres} as a function of M_{max} and $R_{1.6}$. These results are derived from our simulations that use the full 3D EoS tables. As per our notation, green crosses indicate EoSs that incorporate hyperons as degrees of freedom, whereas black plus signs correspond to purely nucleonic EoSs. The figure indicates that it is not possible to clearly distinguish between the two classes of EoSs based on their threshold mass values in this parameter space. The observation confirms the existing conclusion in the literature that these empirical relations hold well for all types of hadronic EoSs, regardless of whether they include exotic degrees of freedom like hyperons. To quantify this universality, we derive a bilinear fit for M_{thres} , according to Eq. (4.13). This fit was performed using data points from the purely nucleonic models only, essentially establishing a benchmark relation expected for matter without exotic components. The goodness of the nucleonic EoSs fit is characterized by an average deviation of approximately $0.03 M_{\odot}$.

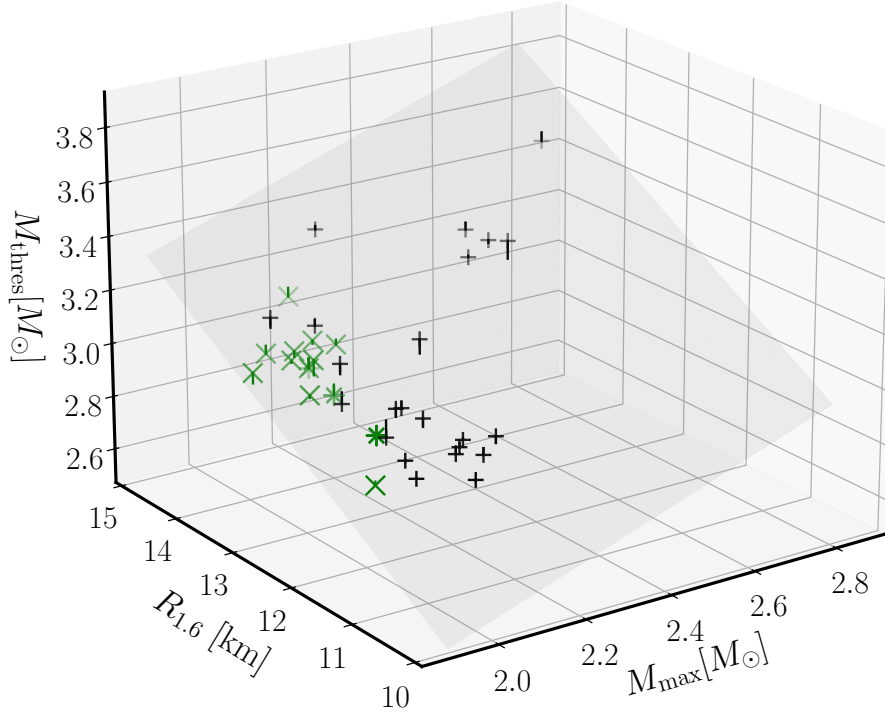


Figure 4.26: Threshold binary mass M_{thres} as a function of M_{max} and $R_{1.6}$. The faded gray plane shows the bilinear fit as described in the text. The small lines suggest the distance of each point to the plane. Purely nucleonic models are indicated by plus signs, models incorporating hyperons are marked with crosses, and models that include both hyperons and Δ baryons are represented by asterisks. The figure was published in [369].

The maximum residual of the fit reads $0.07 M_{\odot}$. When we apply the nucleonic-derived fit to the hyperonic models, which were not included in the fitting procedure, the deviations remain similar. For them, the average deviation from the fit is also approximately $0.03 M_{\odot}$, while the maximum deviation is only marginally lower at $0.06 M_{\odot}$. The fact that the hyperonic models fall so closely onto the relation defined by purely nucleonic EoSs confirms the robust nature of this quasi-universal relation.

The equivalence between the nucleonic and hyperonic classes of EoS models regarding M_{thres} is more evident in Fig. 4.27. This two-dimensional plot displays M_{thres} directly as a function of the linear combination $aM_{\text{max}} + bR_{1.6}$, where a and b are the coefficients derived from the bilinear fit to the purely nucleonic data, as discussed in the preceding paragraph. In this representation, all data points, including those for EoSs that incorporate heavy baryons, are observed to scatter closely around a straight line, which essentially represents the universal hadronic relation projected onto a plane. The tight clustering demonstrates that the presence of hyperons does not alter this particular quasi-universal relation. We have verified that the result holds if other pairs of variables are employed,

such as $(M_{\text{max}}, \Lambda_{1.6})$ or $(M_{\text{max}}, \Lambda_{1.4})$. This universality is an advantage for constraining NS parameters $(M_{\text{max}}, R, \Lambda)$ from observed merger outcomes, as it implies that the derived constraints are largely independent of the presence or absence of hyperons.

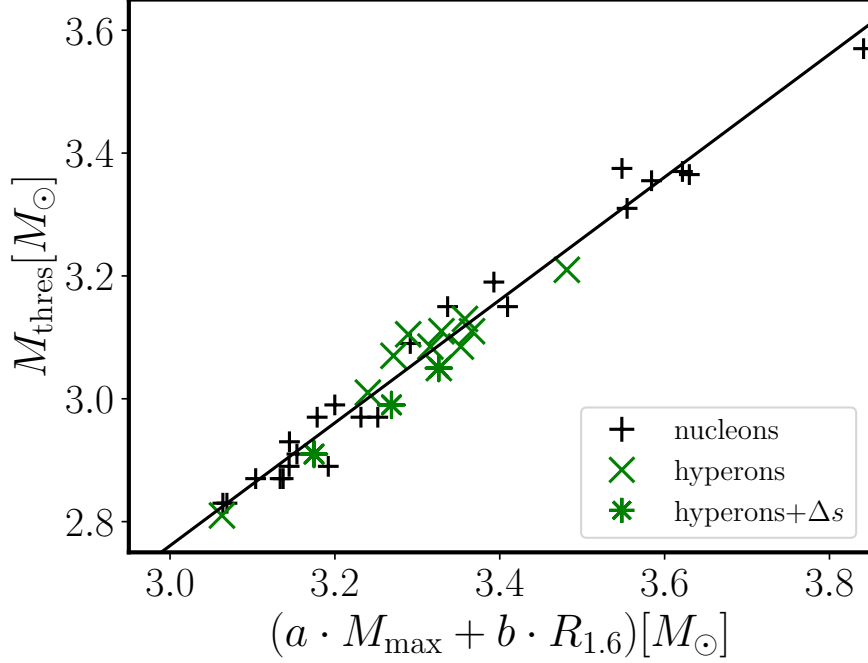


Figure 4.27: 2D plot of M_{thres} as a function of $aM_{\text{max}} + bR_{1.6}$. The coefficients in the linear term $aM_{\text{max}} + bR_{1.6}$ are obtained using the bilinear fit shown in Fig. 4.26. The straight line is the projection of the plane. The figure was published in [369].

It is intriguing that heavy baryons do not exhibit a clearer distinction in these quasi-universal bilinear relations for M_{thres} . As discussed previously, their specific thermal behavior has affected other observables such as the spectral features. To address this query further, we also consider the set of simulations performed with an approximate thermal treatment, adopting a constant thermal index of $\Gamma_{\text{th}} = 1.75$. The results are shown in Fig. 4.28. In this figure, some hyperonic models have small but visible deviations from the bilinear fit. These deviations are characterized by a slightly higher threshold mass, despite assuming the same idealized nucleonic thermal behavior.

Note that the hyperonic and the nucleonic models populate distinct regions of the stellar parameter space, as can be seen in Fig. 4.26. For models incorporating heavy baryons to be realistic and achieve sufficiently high maximum masses for cold NSs (e.g. $M_{\text{max}} \gtrsim 2M_{\odot}$), their low- and intermediate-density regions, where hyperons are still not present, must have a certain stiffness. Consequently, these models often predict relatively large NS radii in the intermediate mass range, even if their maximum mass only slightly exceeds

two solar masses due to softening at higher densities. If these hyperonic models were, hypothetically, purely nucleonic EoSs possessing the same properties (i.e., if their cold EoS was identical to the nucleonic models in these regions of parameter space), our analysis would suggest that the relation is not perfectly bilinear. In the region of the parameter space characterized by a relatively large $R_{1.6}$ combined with relatively small $M_{\max}$, the true relation would actually be slightly increased with respect to a simple bilinear function.

Hence, we conclude that the current nucleonic sample does not properly cover this part of the parameter space. The hyperonic models, when treated with their full temperature-dependent EoS, due to their unique thermal properties, exhibit a lower M_{thres} . This thermal softening effectively compensates for the intrinsic "non-linearity" in the cold EoS parameter space that would otherwise lead to higher M_{thres} values, thus leading to the observed robust bilinear behavior of $M_{\text{thres}}(M_{\max}, R_{1.6})$ when using the full 3D EoS tables. To confirm our conclusions, future studies with a more comprehensive sampling of nucleonic EoSs, particularly exploring these less-covered parameter regimes, would be highly valuable.

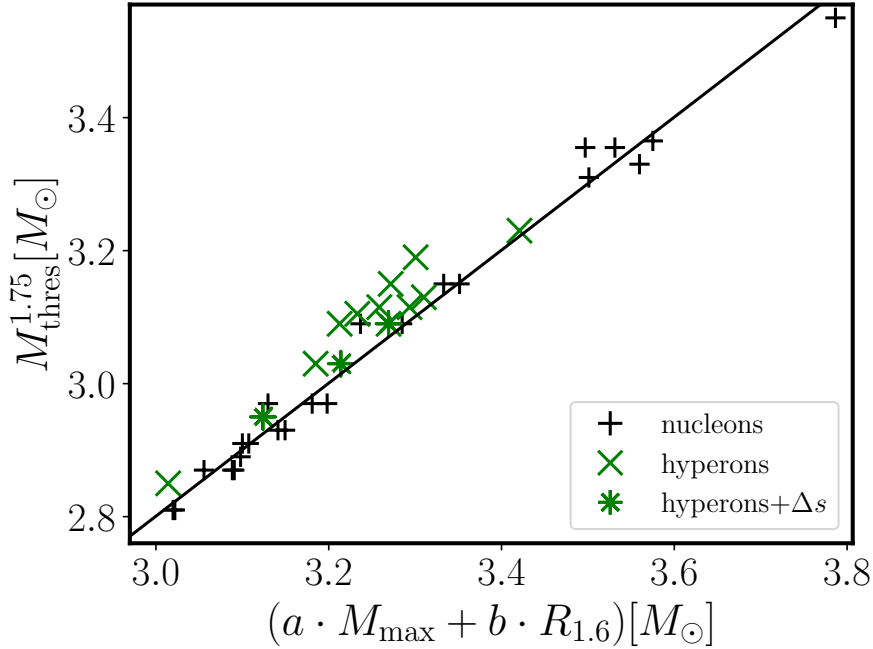


Figure 4.28: Same as Fig. 4.27, but for the $\Gamma_{\text{th}} = 1.75$ run. The figure was published in [369].

The interplay between cold EoS and thermal effects, as revealed in our analysis of M_{thres} , opens a new, albeit challenging, avenue to infer the presence of heavy baryons. Although

the M_{thres} bilinear relations might appear universally valid across hadronic EoSs, another comparison may hint at the subtle yet potentially detectable impact of exotic matter. In Fig. 4.29 we illustrate the difference ΔM_{thres} obtained from the outputs of the two runs for each EoS as a function of M_{thres} . This comparison explicitly shows that the presence of hyperons reduces the threshold mass by roughly $0.05 M_{\odot}$ compared to the reference simulation with an 'idealized' nucleonic thermal behavior. This reduction in M_{thres} for hyperonic models is a direct consequence of the thermal softening and the reduction in $\bar{\Gamma}_{\text{th}}$. A softer EoS, even if only thermally softened, typically leads to a more compact remnant that is more prone to prompt collapse, thus lowering M_{thres} . It is also worth noting that using the simplified $\Gamma_{\text{th}} = 1.75$ approach for purely nucleonic EoSs, one might slightly underestimate the true threshold mass.

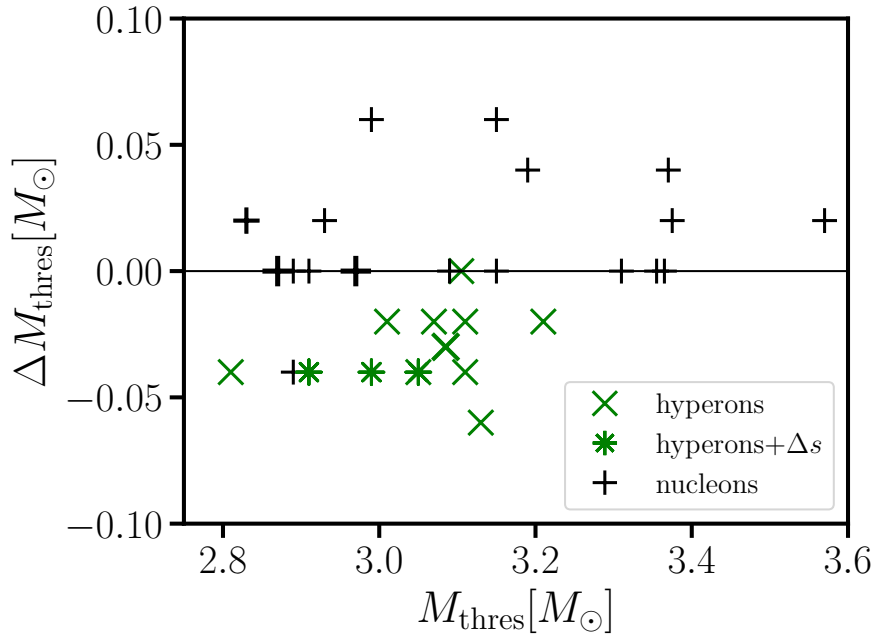


Figure 4.29: ΔM_{thres} as function of M_{thres} . Purely nucleonic models are indicated by plus signs, models incorporating hyperons are marked with crosses, and models that include both hyperons and Δ baryons are represented by asterisks. If two or more models produce equal ΔM_{thres} and M_{thres} , the corresponding data point is labeled with a thicker symbol. The figure was published in [369].

Deducing the presence of hyperons using precise measurements of M_{thres} , however, requires several prerequisites to be fulfilled. Firstly, a highly accurate determination of M_{thres} from an observed GW event is essential. Secondly, such an inference would also demand a precise knowledge of the cold NS EoS, which is crucial for a meaningful comparison of the measured threshold mass with the theoretical expectation for purely

nucleonic EoSs, in a manner analogous to f_{peak} . The deviation from this nucleonic benchmark could then serve as an indicator for the presence of hyperons. We acknowledge, however, that the exact numerical determination of M_{thres} from simulations itself remains subject to various factors, including its precise definition (e.g., the exact collapse criterion), as well as the specific physical model employed, the numerical scheme, and the resolution of the simulation [440, 441].

3.5 Further analysis: Effect of the total mass and the asymmetry in BNS mergers

So far, for simulations that did not suffer prompt collapse, we focused on simulations of symmetric binaries with a total mass of $2.8 M_{\odot}$. We investigated specific effects that hyperons have on the simulation observables. It is important to demonstrate that these effects are not related to the particular initial conditions but are instead a direct consequence of the appearance of heavy baryons in matter. In this subsection, we explicitly show this by focusing our analysis on f_{peak} , extracted from simulations with different initial configurations.

We begin our discussion by analysing how the total mass of the initial binary system can influence the results. We restrict ourselves to the case of symmetric binaries. If the initial mass of the configuration is increased, the average and central densities reached within the remnant increase substantially. The higher density drives the matter into regimes where hyperons may become energetically more favorable.

We present results from a dedicated set of simulations employing heavier binaries. Our study employs the following EoS models: the FSU2H* family, QMC and DNS. This selection is intentional because for the canonical $1.4\text{--}1.4 M_{\odot}$ mergers, those EoSs yielded relatively small or even negative values for Δf_{peak} . For each of these chosen hyperonic EoSs, we perform simulations with binary masses of $2.9 M_{\odot}$, $3.0 M_{\odot}$, or $3.1 M_{\odot}$. The selected configuration is such that it is just below its respective M_{th} for immediate collapse. This selection criterion ensures that the remnants form hyper-massive NSs. Similarly to the canonical simulations, with each EoS we perform two different runs (3D and $\Gamma_{\text{th}} = 1.75$) to obtain the frequency shift, Δf_{peak} .

Fig. 4.30 illustrates the impact that the increased total binary mass has on the dominant post-merger frequency. The results from the more massive binaries are represented by crosses. The plot also contains our canonical merger data, which is represented by dots. For each hyperonic EoS model, a gray line links the more massive data points and their $1.4\text{--}1.4 M_{\odot}$ counterparts.

As anticipated, for more massive and thus more compact remnants, the simulations consistently produce a higher f_{peak} . More significantly, the frequency shift systematically

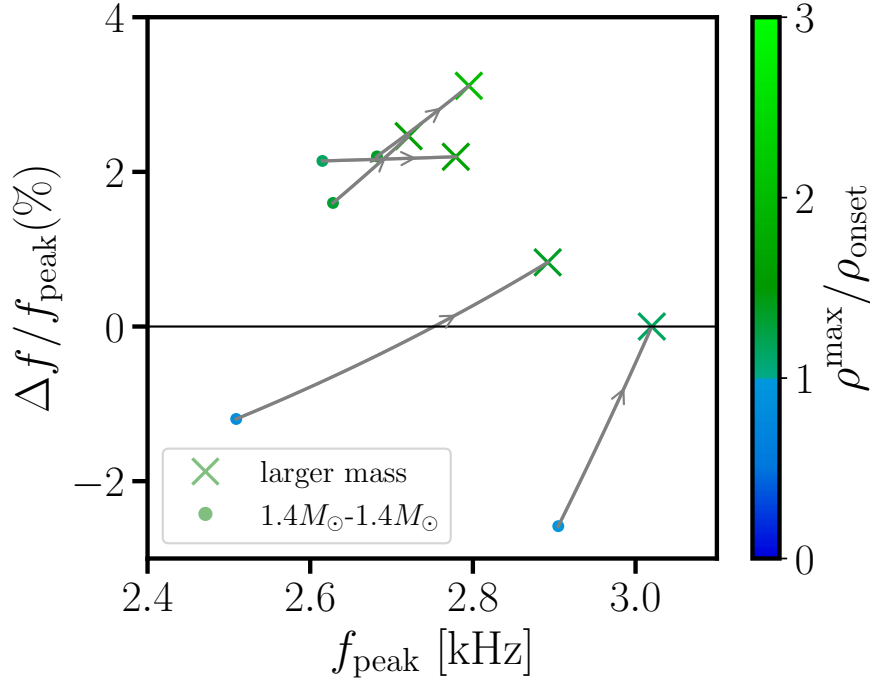


Figure 4.30: Δf_{peak} as a function of f_{peak} in the case of more massive binaries (see the text for detail). For each EoS model, the gray line connects the results of the canonical mass simulations with the higher mass simulations. The arrow indicates the direction towards larger masses. The color bar indicates the ratio between the maximum rest-mass density achieved in the merger remnant within the first 5 milliseconds after the merger and the zero-temperature onset density of heavy baryons ($\rho^{\text{max}}/\rho_{\text{onset}}$) for the more massive configuration. The figure was published in [369].

increases for more massive binaries. The subset of EoS models that initially generated negative Δf_{peak} for the canonical configurations (e.g., DNS, QMC) now show a larger frequency shift. In these specific cases, the additional mass leads to a higher hyperon population compared to the canonical configuration. Nevertheless, even with the increase in mass, the values of the shift for these EoSs are still not as large as for the other models. This behavior is a direct consequence of the abundance of thermal hyperons, which is still relatively low. That is the reason why in Fig. 4.31 $\bar{\Gamma}_{\text{th}}$ remains close to values usually produced by purely nucleonic matter. The anti-correlation between $\bar{\Gamma}_{\text{th}}$ and Δf persists for the more massive binaries as well. This conclusion is also valid for the EoSs that already featured a positive shift for the $1.4\text{--}1.4 M_{\odot}$ binaries simulation, and for which the additional increase in shift for more massive binaries is comparatively smaller.

We now turn our attention to study whether the hyperons imprints on f_{peak} also hold when one simulates asymmetric binaries. We conducted a simulation campaign analogously

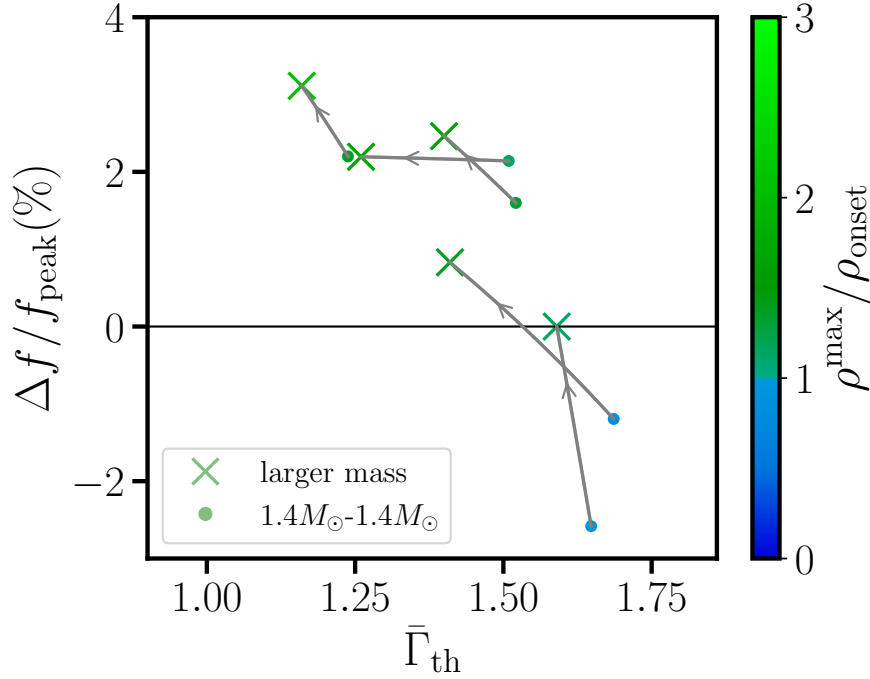


Figure 4.31: $\Delta f/f_{\text{peak}}$ as a function of $\bar{\Gamma}_{\text{th}}$ in the case of more massive binaries (see the text for detail). The notation is the same as in Fig. 4.30. The figure is published in [369].

to the canonic one. With each EoS, with an initial configuration of $M_{\text{tot}} = 2.8 M_{\odot}$ and $M_1/M_2 = 0.8$, two runs of simulations (full 3D EoS and $\Gamma_{\text{th}} = 1.75$) were performed.

Fig. 4.32 shows the relative frequency shift, $\Delta f/f_{\text{peak}}$, from the asymmetric campaign as a function of f_{peak} . A positive frequency shift for the hyperonic EoSs is observed in this case as well. The relative shift in the asymmetric merger can be even more pronounced. For a few hyperonic EoS, the shift is close to 6%, almost 2% higher than the maximum relative shift obtained from our canonical symmetric campaign. The enhanced shift in asymmetric mergers might be attributed to the qualitatively different dynamics and remnant properties of unequal-mass binaries. Asymmetric mergers tend to produce less uniformly rotating remnants, which can lead to larger tidal deformations and potentially different thermalization patterns, which may further enhance the conditions for hyperon production or their impact on the remnant's stiffness.

It is noteworthy that only 13 different hyperonic models are depicted in the figure, even though our calculations took into account 14 EoS models with heavy baryons. The missing data point corresponds to the softest hyperonic EoS model, SFHOY, for which the remnant promptly collapsed in the simulation utilizing the full 3D EoS table. Contrary, using the same EoS in the Γ_{th} run, a hyper-massive NS is formed. We have already

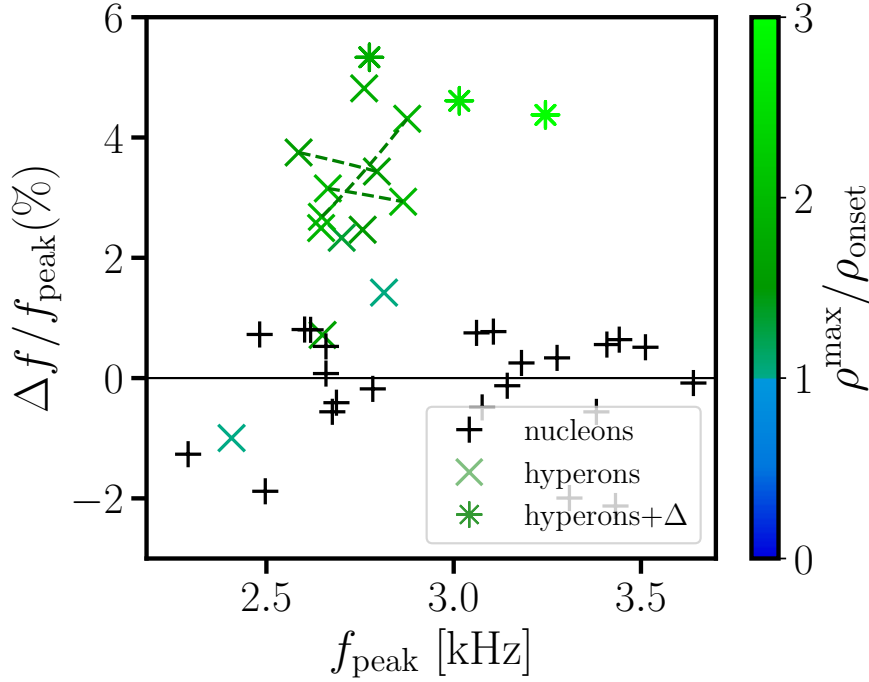


Figure 4.32: $\Delta f_{\text{peak}}/f_{\text{peak}}$ as a function of f_{peak} obtained from the asymmetric campaign. Purely nucleonic models are indicated by plus signs, models incorporating hyperons are marked with crosses, and models that include both hyperons and Δ baryons are represented by asterisks. Three of our models have two dominant peaks present in their GW spectrum, which are roughly equally strong. For those EoSs, both peaks are displayed and connected with a dashed line. The figure was published in [369].

discussed the implications of different thermal behaviors on the threshold mass in the previous subsection. Additionally, the hyperonic EoSs developed in Ref. [258] yield GW spectra with two dominant peaks of comparable strength. In Fig. 4.32, those particular data points are connected with a green dashed line. The precise mechanism responsible for generating these double, near-equally dominant peaks is yet to be resolved.

The relation between $\Delta f/f_{\text{peak}}$ and $\bar{\Gamma}_{\text{th}}$ holds very well in the asymmetric campaign too, as seen in Fig. 4.33. Again, the EoSs that produce a smaller thermal index in the remnant also predict higher frequency shifts.

Consistent with our approach for equal-mass binaries discussed in the beginning of the section, we also extend our investigation to explore the direct quasi-universal relations between Λ and f_{peak} . We have performed a similar procedure to identify the optimal reference mass M for a tight $f_{\text{peak}} - \Lambda_M$ relation for nucleonic models in asymmetric binaries and found that $M = 1.65 M_{\odot}$ is the value for which this relation is tightest.

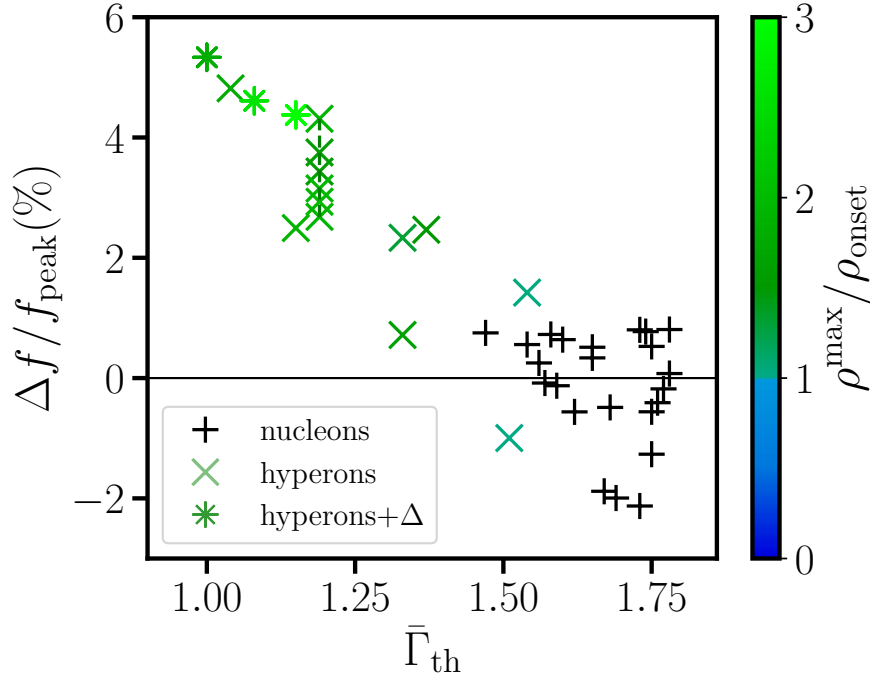


Figure 4.33: $\Delta f/f_{\text{peak}}$ as a function of $\bar{\Gamma}_{\text{th}}$ obtained from the asymmetric campaign. The notation is the same as in Fig. 4.32. The figure was published in [369].

Fig. 4.34 displays the corresponding $f_{\text{peak}} - \Lambda_{1.65}$ relation. Similarly to the equal-mass results, the plot shows a systematically higher frequency for hyperonic models compared to purely nucleonic models for the same value of $\Lambda_{1.65}$. Indeed, the softening effects induced by heavy baryons have a significant influence, irrespective of the initial mass ratio. The qualitative agreement between the equal-mass and $M_1/M_2 = 0.8$ asymmetric binary results strongly suggests that this phenomenon is a general characteristic. We anticipate that for hyperonic EoSs, higher values of f_{peak} , with respect to the nucleonic ones, will be found for binaries with other mass ratios as well. This further reinforces the potential of f_{peak} to serve as an indicator of the hyperon appearance in the dense matter of the remnants.

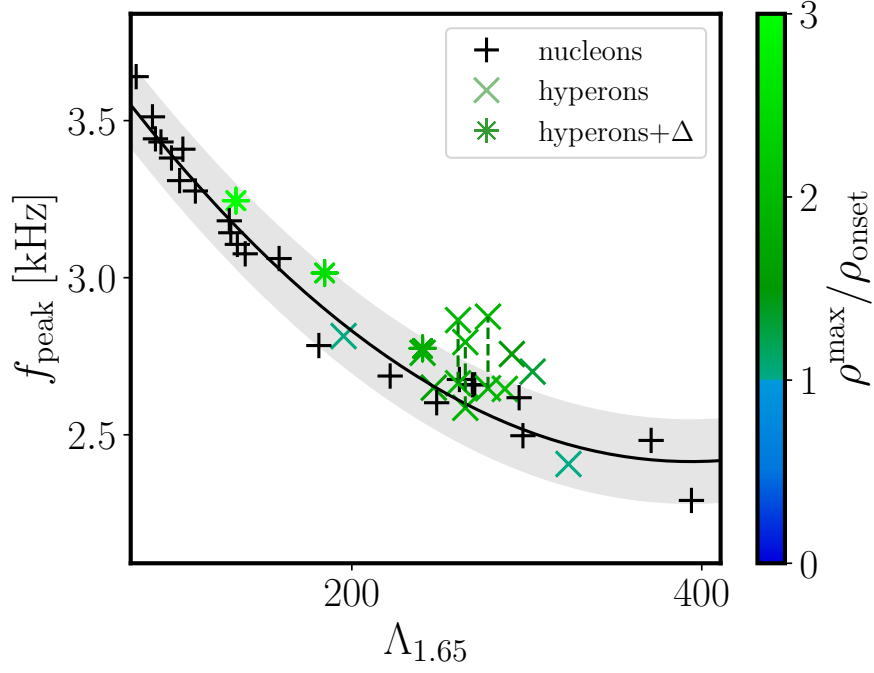


Figure 4.34: f_{peak} as function of $\Lambda_{1.65}$ in asymmetric mergers with a binary mass ratio $M_1/M_2 = 0.8$. Purely nucleonic models are indicated by plus signs, models incorporating hyperons are marked with crosses, and models that include both hyperons and Δ baryons are represented by asterisks. The black curve represents the least-squares quadratic fit to purely nucleonic models. The gray band labels the maximum residual of purely nucleonic models from the fit. Three of our models have two dominant peaks in their GW spectrum, which are roughly equally strong. For those EoSs, both peaks are displayed and connected with a dashed line. The figure was published in [369].

Chapter 5

Summary and Conclusions

The appearance of hyperons in ultra-dense and hot matter is the central topic that has been discussed in this thesis. We approached the topic from two different angles. First, we studied a theoretical formalism in which hyperons can be included at the level of the EoS. Second, we studied the different signatures that they leave on the stellar observables from isolated cold stars and from BNS merger events.

We started the thesis by developing a RMF EoS that explicitly incorporates hyperons as degrees of freedom to describe homogeneous baryonic matter at arbitrary conditions (density, charge fraction, temperature). The choice of the framework was motivated by several factors. As we have explained, RMF models are built upon the principles of relativistic quantum field theory. They inherently preserve relativistic covariance, which is important because in the extreme conditions of neutron star matter, the particles can move at relativistic speeds. Therefore, the models naturally satisfy causality, which can be broken in non-relativistic methods. In addition, RMF models can easily include hyperon degrees of freedom by introducing coupling terms between hyperons and the same mesons that mediate NN interactions (and potentially new ones like the ϕ and σ^* mesons). This allows for a consistent framework to describe the many-body YN and YY interactions.

Our starting point was the FSU2H cold EoS of Refs. [259, 260], which was developed in Chapter 2 taking into account the saturation properties of nuclear matter, some properties of finite nuclei, as well as astrophysics constraints from massive pulsars. We firstly re-examined the model parameters in the hyperonic sector, taking into account the newest results regarding hyperon potentials in symmetric nuclear matter and $\Lambda\Lambda$ bond in Λ matter. Consequently, our updated framework, termed FSU2H*, integrates these improvements. We furthermore extended this model to finite temperature, thereby enabling the calculation

of the EoS for dense and hot matter at various compositions, covering conditions relevant to supernova explosions, PNSs, and BNS mergers.

We dedicated special attention to NS matter in neutrinoless beta-equilibrium, as this scenario is important for the understanding of evolved isolated stars and closely mimics the final phases of PNS evolution. We focused on the results obtained with our nominal FSU2H* model. We again inspected the thermal behavior of the EoS, but this time through the quantity known as the thermal index. Our findings revealed that while the variations in the thermal index with temperature and density are mild for purely nucleonic matter, this is not the case for hypernuclear matter. The appearance of hyperons induces a drop of the thermal index that is stronger at lower temperatures ($T < 20$ MeV) and can even lead to negative values. At higher temperatures, the drop becomes milder but extends over a broader range in density space. These results have significant implications for relativistic simulations of supernovae and BNS mergers, which conventionally rely on a cold EoS supplemented by an approximated thermal behavior, assuming a constant value of the thermal index (often referred to as the Γ -law). Our work demonstrated that this approximation is inadequate to accurately reproduce the true thermal effects, particularly when hyperons become abundant in hot and dense matter conditions encountered in these astrophysical events. To make our results practically useful, we also provided a simple parametrization of the thermal index that can be used in relativistic simulations.

However, we have also extensively discussed that the hyperonic model parameters remain insufficiently constrained because of the substantial uncertainties still persisting in the experimental data for most hyperon potentials. Hence, a part of this thesis was dedicated to systematically exploring how these uncertainties in hyperon potentials influence the hyperonic EoS at finite temperature. We developed two additional parametrizations, FSU2H*L and FSU2H*U, which are based on our nominal FSU2H* model. These parametrizations effectively cover the uncertainty range of the hyperon potentials, representing lower (L) and upper (U) limits, leading to the softest and the stiffest hyperonic EoSs, still consistent with the experimental data. We explicitly showed that the variation of the model propagates to an uncertainty of our predictions both in the composition of matter and the EoS itself. The prediction of the pressure in ultra-dense matter was particularly affected, whereas the entropy per particle and the energy per particle exhibited less sensitivity to variations in the hyperon potentials within matter.

While in Chapter 2 we dealt with the effects that the uncertainties of the hyperonic potentials have on the EoS and the composition of matter, Chapter 3 explained how those effects translate into the predicted global properties of the stars. We focused on the results for the mass-radius relation of the stars, their tidal deformability and the moment of inertia. We also discussed the strangeness factor, as it is closely related with the effects that hyperons induce on the observables.

We first presented calculations at zero temperature in neutrinoless beta-equilibrium and, whenever possible, we compared the obtained results with the constraints from astrophysical measurements. We saw that the hyperonic uncertainties influenced the maximum cold NS mass that the EoS can predict, with a $0.14 M_{\odot}$ difference between the prediction for our stiffest parametrization FSU2H*U and the softest parametrization FSU2H*L. Furthermore, for the most massive stars, the difference of the predicted radius for a star with fixed mass can be up to 10%. However, the impact of the uncertainties on the corresponding radii of maximum mass stars is surprisingly negligible, as it does not exceed 0.2 km. Similar results were obtained for the tidal deformability and for the moment of inertia too. The tidal deformability at high NS masses was especially affected by the uncertainties, as the predictions between our extreme parametrizations can be different by a factor ≈ 2.8 for a star with fixed mass. This fact gives hope that with future measurements of cold star observables, we can significantly constrain the hypernuclear interaction at high densities. Comparing the obtained results with the results for the strangeness number, it was clear that the discrepancy between the parametrizations is the highest when the strangeness number shows a maximum.

We also performed calculations at finite temperature using temperature and lepton profiles for stars motivated by conditions present in the evolution of PNSs both in their early evolution (around 1 s after the explosion of the star) and in late phases (tens of seconds after the explosion of the star). We found that PNSs during their early evolution have lower compactness, and we connected this to the stiffness of the low-density EoS in hot matter. The influence of the hyperonic uncertainty on the observables was analogous to the one at zero temperature. However, it was noticeable that because of the higher lepton fraction in the early evolution, the appearance of hyperons was hindered, leading to a lower strangeness number for those cases compared to the cases where matter is considered to be in neutrinoless beta-equilibrium. Consequently, the impact of the uncertainties is higher on the observables calculated for stars in their late evolution phases.

The specific effects that hyperons induce at the finite-temperature EoS, which were discussed in Chapter 2 and Chapter 3, motivated a deeper investigation that went beyond a single-model description of matter. In Chapter 4 we took 14 different hyperonic EoSs, and in the first part of this chapter we analysed their finite-temperature properties. We confirmed that the drop of the thermal pressure when hyperons appear in matter is present regardless of the model/framework used. We connected this with the appearance of thermal hyperons in matter. In addition, it was shown that hypernuclear matter in general shows a higher heat capacity than ordinary nuclear matter.

The results provoked an extensive simulation campaign that was performed with the aforementioned hyperonic EoSs and also with 23 purely nucleonic EoSs. The results were presented in the second part of Chapter 4. The goal of the campaign was to properly identify possible hyperonic signatures in merger observables and to discuss their effect on

the evolution of different merger quantities. At the same time, the simulation campaign allowed us to study how the thermal behavior in general impacts the outcome of the simulations.

We have observed a small (around a few percent) but systematic increase of the dominant post-merger GW frequency when thermal hyperons are present in remnant matter with respect to cases when they are not present. The frequency shift anti-correlated with the thermal index of matter, confirming the fact that the softening of the finite-temperature EoS induced by thermal hyperons triggers this phenomenon. We corroborate that the characteristic increase did not depend on the initial condition of our simulations. The effect was present in our nominal simulations—symmetric binaries run with a total mass of $2.8M_{\odot}$; in our asymmetric ($M_1/M_2 = 0.8$) binary runs, while keeping the same total mass of $2.8M_{\odot}$; and in our more massive symmetric runs.

We also explored the influence of hyperons on other features in the GW spectra. We explicitly showed that the secondary peak, f_{2-0} , suffers a similar increase when thermal hyperons are present in the matter, while the secondary peak, f_{spiral} , appears not to be affected by the appearance of hyperons. We related the different qualitative behaviors of the two features with the different mechanisms by which they are produced.

The average temperature in the remnants was also studied. It was clear that the higher specific heat of hypernuclear matter induced lower temperatures in the remnants compared to the cases where only nuclear matter was present.

We also studied the dynamical mass ejecta and the threshold mass for prompt black hole collapse. When hyperonic EoSs are used, the dynamical mass ejected appears to be greater than the one obtained with nucleonic EoSs, for cases where both EoSs yield similar cold NS properties. Furthermore, hyperonic EoSs lead to a lower threshold mass for prompt black hole formation (by approximately $0.05M_{\odot}$) compared to their nucleonic counterparts, assuming the same cold NS stellar parameters.

A word of caution is also needed. Some of the hyperon effects discussed in the thesis may be mimicked by other forms of exotic matter, like quark matter [405, 404] or even nuclear matter [406, 407, 408]. In Appendix E we address part of this issue by developing a toy model that incorporates hyperons in nuclear matter that already has a particular low thermal index. Our preliminary results show that even if nucleonic matter produces a particularly low thermal index, the appearance of hyperons triggers an additional thermal softening. Furthermore, in some recent works [220, 443] it was shown that explicitly momentum-dependent hypernuclear potentials can significantly alter the properties of the cold EoS. Future work should address whether the conclusions obtained in the thesis still hold when a more complex, momentum-dependent interaction is employed for obtaining the hyperonic finite-temperature EoS.

The results presented in the thesis represent a small but important step toward discovering the true composition of NSs. Besides our work, many other authors have also tried to find hyperon signatures in other observables, both at zero temperature and at finite temperature [444, 445, 446, 447, 448, 449, 450, 451, 452]. Identifying clear hyperon signatures both at the EoS level and at the level of observables will allow future observations to significantly constrain the EoS of hot and dense matter. On the one hand, our results might motivate the use of new experimental techniques such as femtoscopy [453, 454, 455, 456, 457, 458, 459, 460, 461, 462, 463, 464, 465, 466, 467] to further constrain the uncertainty that persists in our understanding of the YN and YY interactions. On the other hand, they highlight the need for advanced observatories for multi-messenger astrophysics that would give hints about the composition of ultra-dense matter [400, 468, 469, 470, 471, 472]. The input of all different fields will finally close the "hyperon puzzle," hopefully confirming the presence of hyperons in the dense matter of neutron stars.

Appendix A

Expansion of the cold nuclear equation of state

We have explained in Chapter 2 that the core of isolated cold stars is made of infinite homogeneous matter. At densities close to the saturation density, the appearance of hyperons is unlikely, so nucleons are the only baryons present in matter. The interaction between nucleons dictates the behavior of matter in those conditions. As the EoS can be obtained in different frameworks, it is important to be able to characterize it with some useful nuclear-matter properties. In this appendix, we explain how one can expand the EoS in powers of density and asymmetry with parameters that can be extracted from nuclear physics experiments.

When a cold infinite asymmetric nuclear matter is considered, the EoS is normally expressed through the energy per particle of the system:

$$E/A(\rho_B, \beta) = \varepsilon_B(\rho_B, \beta)/\rho_B - m_N, \quad (\text{A.1})$$

where $\rho_B = \rho_p + \rho_n$ is the baryon density, with ρ_p and ρ_n the proton and neutron number densities respectively; m_N is the mass of the nucleons and $\beta = (\rho_n - \rho_p)/\rho_B$ is the isospin asymmetry. As in Chapter 2, with ε_B we label the energy density of the system. If one expands E/A around β , the following relation can be obtained:

$$E/A(\rho_B, \beta) = E/A_{\text{SNM}}(\rho_B) + \beta^2 S_2(\rho_B) + \dots + . \quad (\text{A.2})$$

We will restrict to the second order in the expansion only. The terms proportional to the odd powers of β are not present in Eq. (A.2) due to the symmetry of the strong force that acts between proton-proton and neutron-neutron pairs. We can identify the first term of the expression, $E/A_{\text{SNM}}(\rho_B)$ as the contribution that comes from SNM.

The second coefficient in the expression is the so-called symmetry energy, $S_2(\rho_B)$. If $E/A_{\text{PNM}}(\rho_B) = E/A(\rho_B, \beta = 1)$ labels the energy per particle in PNM, then up to second order the following relation holds

$$S_2(\rho_B) = E/A_{\text{PNM}}(\rho_B) - E/A_{\text{SNM}}(\rho_B). \quad (\text{A.3})$$

We can see that the symmetry energy, in fact, gives the energy that is necessary for converting protons from SNM to neutrons in PNM.

The majority of the nuclear physics experiments test the cold part of the EoS at densities close to the saturation one in SNM. Since this density corresponds to the global minimum of E/A_{SNM} , it holds

$$\left. \frac{\partial E/A_{\text{SNM}}}{\partial \rho_B} \right|_{\rho_B=\rho_0} = 0. \quad (\text{A.4})$$

By introducing the dimensionless parameter:

$$x = \frac{\rho_B - \rho_0}{3\rho_0}, \quad (\text{A.5})$$

the energy per particle of SNM can be expanded around saturation density as

$$E/A_{\text{SNM}} = E_0 + \frac{1}{2}K_0x^2 + \dots + . \quad (\text{A.6})$$

We recognize that the first term represents the value of the energy per particle of SNM at saturation density (e.g., the volume term in the liquid drop model description of the nuclei [473]). The second term K_0 is known as the incompressibility coefficient

$$K_0 = 9\rho_0 \left(\frac{\partial^2 E/A_{\text{SNM}}}{\partial \rho_B^2} \right) \bigg|_{\rho_B=\rho_0}. \quad (\text{A.7})$$

This coefficient is most directly probed by the iso-scalar giant monopole resonance, often referred to as the "breathing mode" of the nucleus [474, 475, 476].

The expansion that we wrote in Eq. (A.2) is usually referred to as parabolic approximation. Same as for the E/A_{SNM} , the symmetry energy, S_2 , can be also expand around saturation density, or equivalently, around the dimensionless parameter x

$$S_2(x) = E_{\text{sym},0} + L_0x + \frac{1}{2}K_{\text{sym},0}x^2. \quad (\text{A.8})$$

The parameter $E_{\text{sym},0}$ labels the value, L_0 represents the slope, and $K_{\text{sym},0}$ gives the curvature of the symmetry energy, all of them computed at saturation density. In this expansion, even the linear term is present. Since the pressure and the energy per particle

are related by the expression

$$P = \rho_B^2 \left(\frac{\partial E/A}{\partial \rho_B} \right), \quad (\text{A.9})$$

we can relate L_0 to the pressure of PNM at saturation density P_{PNM} :

$$L_0 = 3\rho_0 \left(\frac{\partial S_2}{\partial \rho_B} \right) \Big|_{\rho_B=\rho_0} = \frac{3P_{\text{PNM}}}{\rho_0}. \quad (\text{A.10})$$

The pressure P_{PNM} , together with $E_{\text{sym},0}$, can be extracted from the measurements of the heavy nuclei neutron-skin thickness and from elliptic flow in collisions of neutron-rich heavy-ion systems at intermediate energies [477, 478, 479, 480, 481]. The curvature of the symmetry energy

$$K_{\text{sym},0} = 9\rho_0 \left(\frac{\partial^2 S_2}{\partial \rho_B^2} \right) \Big|_{\rho_B=\rho_0}, \quad (\text{A.11})$$

can be extracted in a similar fashion like L_0 and $E_{\text{sym},0}$. However, its effects are generally smaller and more subtle than those of L_0 and $E_{\text{sym},0}$.

Finally, combining the expressions (A.2), (A.6) and (A.8), the energy per particle at arbitrary asymmetry and density can be written as

$$E/A(x, \beta) = E_0 + E_{\text{sym},0}\beta^2 + L_0x\beta^2 + \frac{1}{2} (K_0 + \beta^2 K_{\text{sym},0}) x^2 + \dots +. \quad (\text{A.12})$$

For a more elaborate discussion about the parametrization of the cold EoS, see Refs. [279, 482, 483, 484, 485]. For recently developed software that allows for comparing different theoretical predictions of dense matter properties and the interplay with nuclear physics and astrophysics data, see Ref. [486].

Appendix B

Interplay of hyperon potentials

In this appendix, we give an expanded discussion on the treatment of hyperonic potential uncertainties within the FSU2H* hadronic model, particularly focusing on the justification of the use of FSU2H*L and FSU2H*U models. As we have discussed in the main part of the thesis, the hyperon effective interactions, due to the scarcity of experimental hypernuclear data, suffer from considerable uncertainties. That is reflected in the model parameters that govern hyperon-meson couplings. The FSU2H*L (Lower limit) and FSU2H*U (Upper limit) models were developed to represent the extreme bounds of the accepted uncertainty ranges for hyperonic potentials. However, some of the quantities can be sensitive to the interplay between the values of the hyperon potentials, since we are dealing with a multidimensional parameter space. Although in principle one could build different parametrizations that address this issue properly, their use is impractical as relativistic simulations are computationally expensive. Hence, we want to check whether our two extreme models cover a wide range of uncertainty space, which would make them suitable for use in relativistic simulations.

We built 18 different additional parametrizations that specifically test the interplay between the hyperon potentials. The new parametrizations are designed to favor/disfavor a particular hyperon species by setting the correspond hyperon potential to its lowest/highest value, which is still compatible with the hypernuclear experimental constraints. At the same time, we systematically change the $\Lambda\Lambda$ bond energy that the parametrizations predict. Three different values are taken into account 0, -0.67, -6.0 MeV, which are responsible for the strength of the hyperon-hyperon interactions. With the obtained parametrizations, several calculations are performed. We inspect the composition and the pressure of neutrinoless β -stable matter for cold matter and for matter at constant entropy per baryon $S/A = 2$. In this way, we test how FSU2H*L and FSU2H*U models cover the uncertainty of the EoS at zero and finite temperature.

We start by examining the composition of cold β -stable matter computed with our additional parametrizations and our two extreme models (Fig. B.1). The shaded areas in the graph show the possible range of how much of each particle type is present at different densities. These ranges are defined by the solid lines, which represent the highest and lowest amounts of each particle observed across all tested conditions. Hence, the shaded regions illustrate the uncertainty in a particle's abundance. The predictions from the FSU2H*L and FSU2H*U models are shown as dashed and dash-dotted lines, respectively. It is clear that the specific hyperon abundance is heavily influenced by the interplay of the uncertainties in hyperon interactions. The extreme models only partially cover the possible range for hyperons, particularly for the Σ^- and Ξ^- hyperons, where a significant uncertainty in the potentials exists. The conclusion still holds at finite

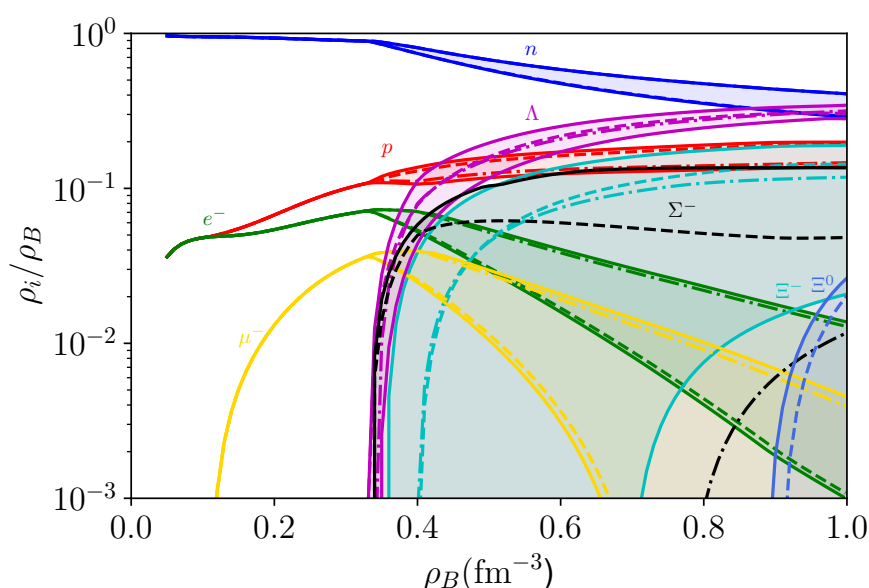


Figure B.1: Composition of cold β -stable matter computed with the additional 18 parametrizations (shaded regions) and with our two extreme models (dashed lines - FSU2H*L; dashed-dotted lines - FSU2H*U). The figure is published in [289].

temperature, as one can see from Fig. B.2, where the composition is presented for β stable matter at constant entropy per particle of $S/A = 2$. However, if one focuses on the total abundance of non-hyperonic species, it is noticeable that the FSU2H*L and FSU2H*U parametrizations actually do describe the uncertainties very well, both at zero and finite temperature, as the abundance curves predicted by these models lie at the edge of the uncertainty regions. As the total hyperonic content is tied to the non-hyperonic one, one can conclude that the uncertainty of the *total* number of hyperons in matter is well described with the two extreme models.

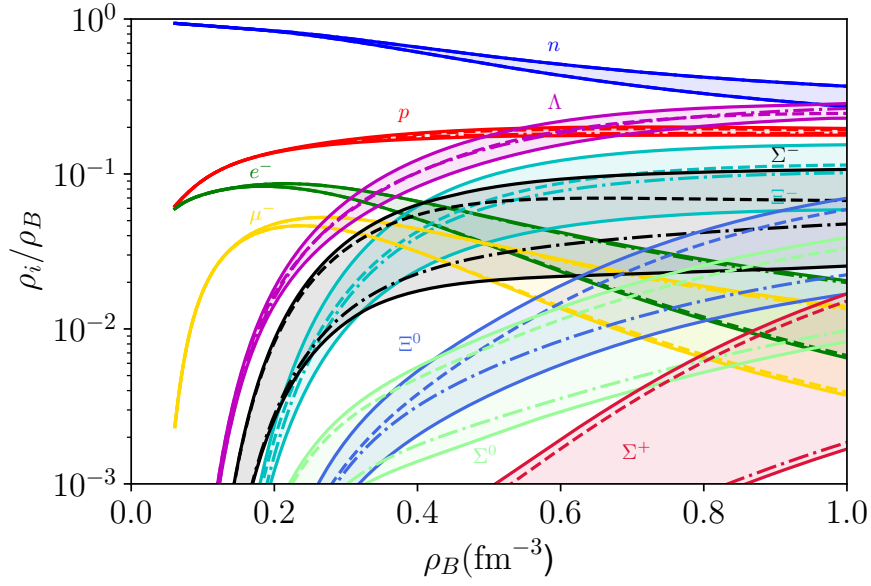


Figure B.2: Composition of β -stable matter at $S/A = 2$ computed with the additional 18 parametrizations (shaded regions) and with our two extreme models (dashed lines - FSU2H*L; dashed-dotted lines - FSU2H*U). The figure is published in [289].

While the composition is sensitive to the interplay between the hyperonic potentials, the same does not hold for the pressure of matter. In Fig. B.3 and Fig. B.4 the β -stable matter pressure as a function of density is presented, for cold matter and matter at $S/A = 2$, respectively. The results obtained with the additional 18 parametrizations are marked with dashed lines. The results with the FSU2H*L model (blue) and the FSU2H*U model (red) are labeled with solid lines. We can see that in fact the two models FSU2H*L and FSU2H*U control exceptionally well the possible stiffness of the EoS. The results should not surprise us. As we already saw in the composition plots, the uncertainty of the total amount of hyperons present in matter is described well with our two extreme models. The reduction of the degeneracy (e.g. the effect that hyperons have on the stiffness of the EoS) does not depend on the specific hyperon present in matter but rather on the amount of hyperons that are produced.

In Chapter 3, using our extreme models, we also studied the global properties of NSs. Therefore, it is also of special interest to test whether our models cover well the uncertainty that arises in the description of global properties of stars from the uncertainty in the interaction. To that end, the mass-radius relation and tidal deformability have been computed with the additional 18 parametrizations and the results are compared with the calculations obtained from the FSU2H*L and the FSU2H*U model. In the same way as for the composition and for the pressure, calculations for two conditions are computed:

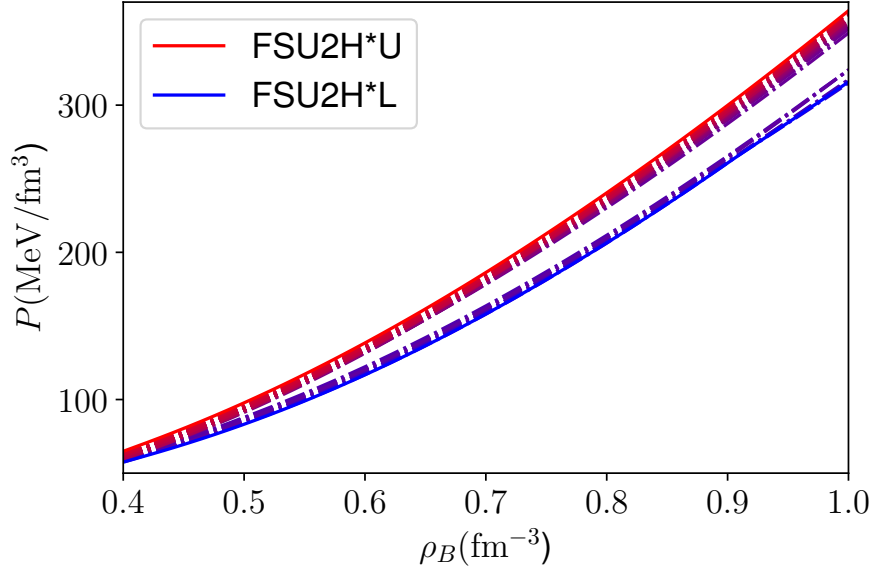


Figure B.3: Pressure of cold β -stable matter computed with the additional 18 parametrizations (dashed lines) and with our two extreme models (solid blue line - FSU2H*L and solid red line - FSU2H*U). The figure is published in [289].

β -stable cold stars and β -stable stars at $S/A = 2$. The mass-radius relation results are shown in Fig. B.5 for cold star matter, and in Fig. B.6 for hot star matter. Fig. B.7 and Fig. B.8 represent the tidal deformability results for cold and hot stars, respectively. We keep the notation from the previous plots and the results obtained with the additional 18 parametrizations are marked with dashed lines, while the ones with the FSU2H*L model (blue) and the FSU2H*U model (red) are labeled with solid lines. In all of the shown plots, the results from the additional parametrizations lie in between the predictions calculated with our extreme models. Indeed, regardless of the temperature, the FSU2H*L and the FSU2H*U models cover well the uncertainty in the global properties of the stars.

We conclude that the FSU2H*L and FSU2H*U models can serve as valuable tools for quantifying uncertainties in our framework that arise from the lack of precise data that constrain the YN and YY interactions. These two extreme parametrizations indeed encompass the broad spectrum of uncertainties associated with the total hyperonic content within dense matter. By exploring the lower (FSU2H*L) and upper (FSU2H*U) limits of hyperonic potentials, these models effectively cover the potential impact of hyperons on the thermodynamic quantities that govern neutron star properties, both at zero and finite temperature. Although these models may not fully resolve the variability of each hyperonic species, their primary strength lies in capturing the uncertainty band arising from the collective hyperonic presence. Therefore, in the absence of specific

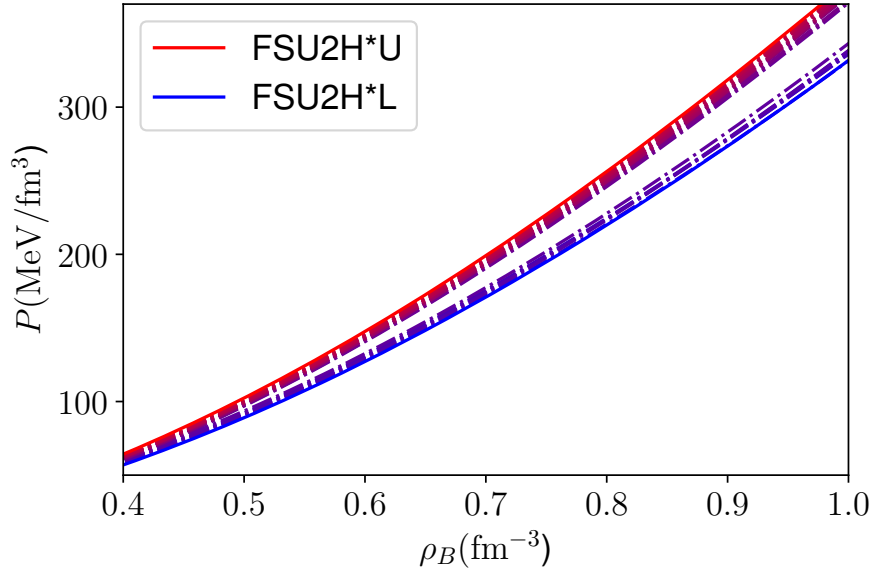


Figure B.4: Pressure of β -stable matter at $S/A = 2$ computed with the additional 18 parametrizations (dashed lines) and with our two extreme models (solid blue line - FSU2H*L and solid red line - FSU2H*U). The figure is published in [289].

observational constraints that are uniquely sensitive to the exact abundance of individual hyperonic species, the FSU2H*L and FSU2H*U models can be employed in relativistic simulations of dynamic astrophysical phenomena to study how the uncertainty propagates to their observables/quantities. This includes, but is not limited to the modeling of BNS mergers and core-collapse supernovae.

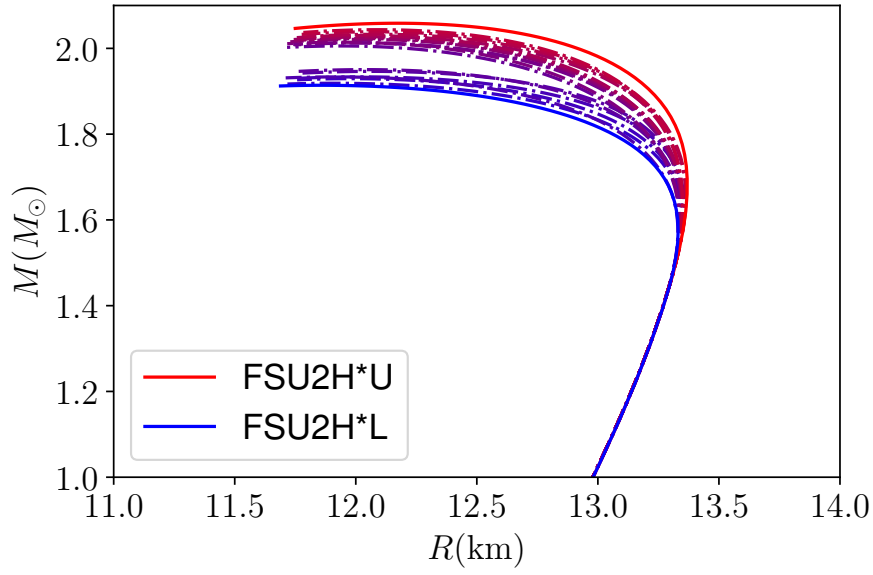


Figure B.5: $M(R)$ relation of cold β -stable stars computed with the additional 18 parametrizations (dashed lines) and with our two extreme models (solid blue line - FSU2H*L and solid red line - FSU2H*U). The figure is published in [289].

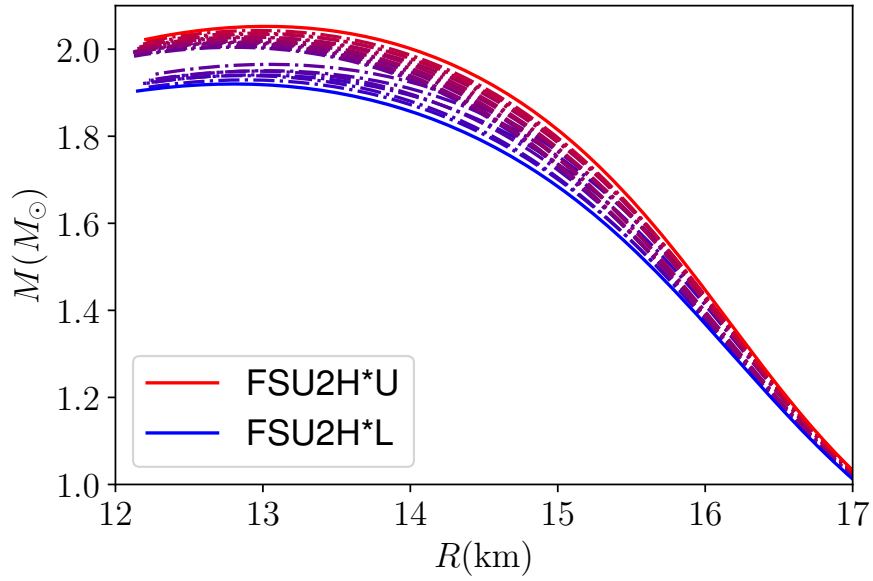


Figure B.6: $M(R)$ relation of hot β -stable stars at $S/A = 2$ computed with the additional 18 parametrizations (dashed lines) and with our two extreme models (solid blue line - FSU2H*L and solid red line - FSU2H*U). The figure is published in [289].

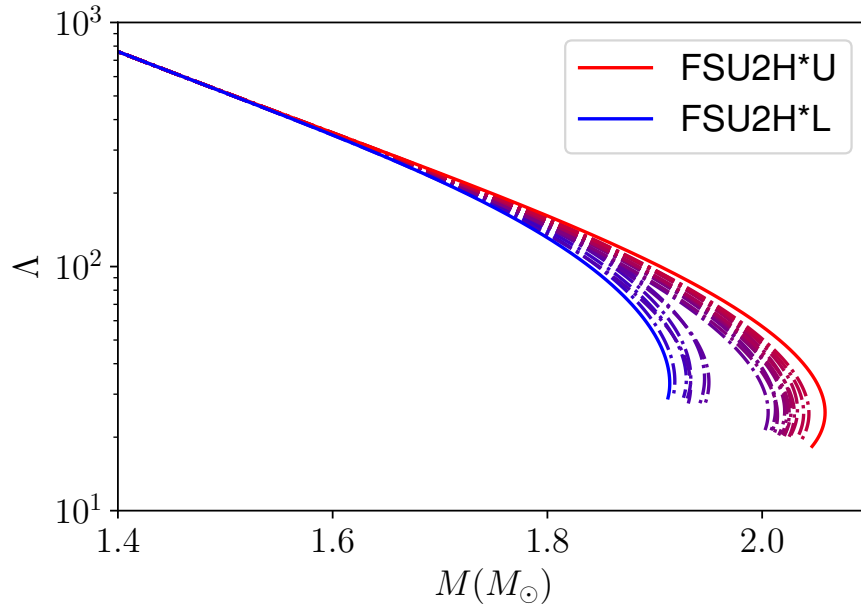


Figure B.7: $\Lambda(M)$ relation of cold β -stable stars computed with the additional 18 parametrizations (dashed lines) and with our two extreme models (solid blue line - FSU2H*L and solid red line - FSU2H*U). The figure is published in [289].

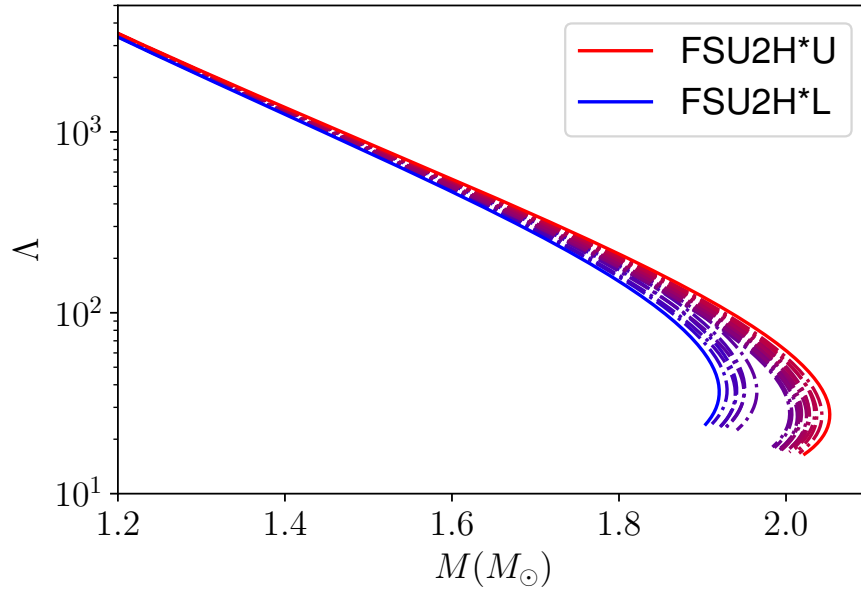


Figure B.8: $\Lambda(M)$ relation of hot β -stable stars at $S/A = 2$ computed with the additional 18 parametrizations (dashed lines) and with our two extreme models (solid blue line - FSU2H*L and solid red line - FSU2H*U). The figure is published in [289].

Appendix C

Cold star properties of different EoSs and simulation results

This appendix lists some properties of cold neutron stars, calculated using the various EoSs used in Chapter 4. It also presents selected results from the simulation campaign described in the same chapter. Table C.1 shows the properties of the cold stars and the results of the $1.4M_{\odot} - 1.4M_{\odot}$ simulation campaign obtained with the hyperonic EoSs. Table C.2 again shows the properties of cold stars and the results of the $1.4M_{\odot} - 1.4M_{\odot}$ simulation campaign, but for nucleonic EoSs utilized in our work. In Table C.3 the results of the merging simulations of more massive NSs are given. Finally, Table C.4 presents the results of our asymmetric simulation campaign ($M_{\text{tot}} = 2.8M_{\odot}$, $q = 0.8$).

EoS	M_{max} [$M_{\odot}$]	$R_{1.4}$ [km]	f_{peak} [kHz]	$f_{\text{peak}}^{1.75}$ [kHz]	f_{2-0} [kHz]	$f_{2-0}^{1.75}$ [kHz]	f_{spiral} [kHz]	$f_{\text{spiral}}^{1.75}$ [kHz]	M_{ej} [$10^{-3}M_{\odot}$]	$\bar{\Gamma}_{\text{th}}$	M_{thresh} [$M_{\odot}$]	$M_{\text{thresh}}^{1.75}$ [$M_{\odot}$]	Ref.
BHBA ϕ	2.10	13.21	2.76	2.68	1.67	1.64	1.98	1.95	4.7	1.37	3.11	3.13	[272]
DD2Y	2.03	13.21	2.82	2.73	1.83	1.68	2.07	2.11	5.4	1.08	3.05	3.09	[273]
DDLS(30)-Y	2.02	12.87	2.69	2.63	1.62	1.57	1.97	2.03	3.7	1.28	3.11	3.11	[258]
DDLS(50)-Y	2.00	13.19	2.74	2.67	1.74	1.60	2.02	2.00	3.6	1.29	3.09	3.12	[258]
DDLS(70)-Y	1.98	13.56	2.71	2.64	1.73	1.61	1.97	1.93	3.3	1.29	3.09	3.12	[258]
DNS	2.09	14.04	2.51	2.54	1.49	1.50	1.78	1.81	3.0	1.69	3.21	3.23	[368]
FSU2H*	2.01	13.18	2.63	2.59	1.55	1.49	2.03	1.99	5.6	1.52	3.11	3.15	[289]
FSU2H*L	1.91	13.16	2.68	2.62	1.63	1.51	1.93	1.89	6.7	1.24	3.07	3.09	[309]
FSU2H*U	2.06	13.17	2.62	2.56	1.52	1.44	1.91	1.89	7.9	1.51	3.13	3.19	[309]
QMC-A	1.99	12.89	2.91	2.98	1.85	1.85	2.17	2.22	6.0	1.65	3.01	3.03	[256]
DD2Y Δ 1.1-1.1	2.04	12.96	3.03	2.93	2.07	1.95	2.26	2.24	7.1	1.08	2.99	3.03	[251]
DD2Y Δ 1.2-1.1	2.05	12.27	3.26	3.14	2.09	1.99	2.48	2.49	10.0	1.18	2.91	2.95	[251]
DD2Y Δ 1.2-1.3	2.03	13.21	2.82	2.72	1.81	1.69	2.04	2.08	4.0	0.99	3.05	3.09	[251]
SFHOY	1.99	11.89	3.60	3.46	2.65	2.52	2.94	2.69	33.5	1.38	2.81	2.85	[248]

Table C.1: Hyperonic EoSs used to simulate equal-mass binary mergers a total mass of $2.8 M_{\odot}$. Columns give the acronyms of the EoS models, maximum mass of cold non-rotating NS, dominant postmerger frequency using a full 3D EoS, dominant postmerger frequency using constant thermal index $\Gamma_{\text{th}} = 1.75$, secondary peak f_{2-0} using a full 3D EoS, secondary peak f_{2-0} using a constant thermal index $\Gamma_{\text{th}} = 1.75$, secondary peak f_{spiral} using a full 3D EoS, secondary peak f_{spiral} using a constant thermal index $\Gamma_{\text{th}} = 1.75$, dynamical mass ejected in the first 10 ms after merging, mass-averaged thermal index of the remnant's fluid in a time windows of 5 ms, starting 2.5 ms after merging, threshold mass obtained using the full temperature dependent EoS, threshold mass obtained using constant thermal index $\Gamma_{\text{th}} = 1.75$ and references of the original works where the models are presented. The table has been published in Ref. [369].

EoS	$M_{\max}$ [$M_{\odot}$]	$R_{1.4}$ [km]	f_{peak} [kHz]	$f_{\text{peak}}^{1.75}$ [kHz]	f_{2-0} [kHz]	$f_{2-0}^{1.75}$ [kHz]	f_{spiral} [kHz]	$f_{\text{spiral}}^{1.75}$ [kHz]	M_{ej} [$10^{-3}M_{\odot}$]	$\bar{\Gamma}_{\text{th}}$	M_{thresh} [$M_{\odot}$]	$M_{\text{thresh}}^{1.75}$ [$M_{\odot}$]	Ref.
APR	2.20	11.57	3.51	3.46	2.26	2.26	2.79	2.71	17.5	1.74	2.87	2.87	[378, 379]
DD2	2.42	13.22	2.64	2.68	1.66	1.62	2.01	2.01	4.7	1.78	3.31	3.31	[338, 271]
DD2F	2.08	12.40	3.30	3.30	2.30	2.27	2.50	2.44	10.1	1.66	2.89	2.93	[338, 380]
DDLS(30)-N	2.48	12.87	2.62	2.62	1.48	1.45	1.95	1.94	3.6	1.83	3.38	3.36	[258]
DDLS(50)-N	2.47	13.19	2.56	2.60	1.40	1.47	1.99	2.06	3.9	1.78	3.36	3.36	[258]
DDLS(70)-N	2.46	13.56	2.62	2.60	1.58	1.52	1.91	1.89	2.4	1.79	3.37	3.37	[258]
DSH Fiducial	2.17	11.73	3.44	3.40	2.26	2.26	2.55	2.67	16.9	1.77	2.89	2.89	[389]
DSH Large Mmax	2.22	12.65	2.93	2.91	1.74	1.76	2.18	2.23	7.2	1.79	3.15	3.09	[389]
DSH Large SL	2.16	11.76	3.51	3.46	2.41	2.30	2.72	2.72	17.4	1.52	2.87	2.87	[389]
DSH Large R	2.13	12.44	3.16	3.18	2.10	2.10	2.46	2.41	12.7	1.72	2.97	2.97	[389]
DSH Small SL	2.18	11.70	3.31	3.33	2.07	2.16	2.52	2.54	13.6	1.76	2.91	2.91	[389]
DSH Smaller R	2.14	11.29	3.62	3.60	2.40	2.50	3.04	2.96	26.2	1.72	2.83	2.81	[389]
FSU2R	2.06	12.87	2.80	2.81	1.71	1.71	2.05	2.13	5.5	1.81	3.09	3.09	[259]
FTNS	2.22	11.46	3.34	3.40	2.10	2.12	2.49	2.59	15.6	1.73	2.93	2.91	[333, 381]
GS2	2.09	13.60	2.73	2.70	1.68	1.71	2.01	1.98	3.0	1.76	3.15	3.15	[336]
LPB	2.10	12.37	3.23	3.23	2.19	2.18	2.43	2.39	8.1	1.68	2.99	2.93	[382, 383]
LS220	2.04	12.96	3.09	3.06	2.08	2.10	2.27	2.25	5.0	1.54	2.97	2.97	[156]
LS375	2.71	13.95	2.44	2.44	1.39	1.40	1.77	1.74	2.0	1.63	3.57	3.55	[156]
SFHo	2.06	11.89	3.43	3.45	2.50	2.38	2.77	2.71	20.5	1.62	2.87	2.87	[271, 370]
SFHx	2.13	11.98	3.16	3.18	2.03	2.01	2.39	2.37	11.0	1.82	2.97	2.97	[271, 370]
SRO(SLy4)	2.05	11.72	3.51	3.50	2.44	2.48	3.00	2.96	10.0	1.78	2.83	2.81	[384, 385]
TM1	2.21	14.47	2.38	2.40	1.31	1.36	1.66	1.71	1.7	1.82	3.37	3.33	[384, 385]
TMA	2.01	13.79	2.58	2.57	1.54	1.57	1.78	1.73	2.1	1.74	3.19	3.15	[388, 387]

Table C.2: Same as Table C.1, but for nucleonic EoSs. The table has been published in Ref. [369].

EoS	Simulated binaries	f_{peak} [kHz]	$f_{\text{peak}}^{1.75}$ [kHz]	$\bar{\Gamma}_{\text{th}}$	$f_{\text{peak}}^{1.4M_{\odot}}$ [kHz]	$f_{\text{peak}}^{1.75, 1.4M_{\odot}}$ [kHz]	$\bar{\Gamma}_{\text{th}}^{1.4M_{\odot}}$	Ref.
DNS	$1.55M_{\odot} - 1.55M_{\odot}$	2.89	2.87	1.41	2.51	2.54	1.69	[368]
FSU2H*	$1.45M_{\odot} - 1.45M_{\odot}$	2.71	2.65	1.40	2.63	2.59	1.52	[289]
FSU2H*L	$1.45M_{\odot} - 1.45M_{\odot}$	2.80	2.71	1.16	2.68	2.62	1.24	[309]
FSU2H*U	$1.50M_{\odot} - 1.50M_{\odot}$	2.78	2.72	1.26	2.62	2.56	1.51	[309]
QMC-A	$1.45M_{\odot} - 1.45M_{\odot}$	3.02	3.02	1.59	2.91	2.98	1.65	[256]

Table C.3: Sample of hyperonic EoSs that we have used to simulate more massive equal-mass binaries. Third to fifth column provide the dominant frequency peak obtained using 3D EoS tables, dominant frequency peak obtained using the approximate thermal treatment and the mass and time averaged thermal index in the remnant. As a reference, sixth to eighth column provide the same quantities but for the less massive, $1.4M_{\odot} - 1.4M_{\odot}$ merger simulations. The ninth column gives references of the original works where the models are presented. The table has been published in Ref. [369].

EoS	Species	f_{peak}	$f_{\text{peak}}^{1.75}$	ρ_{max}/ρ_0	$\bar{\Gamma}_{\text{th}}$	$\Lambda_{1.65}$	Ref.
BHBA ϕ	N-Y	2.65	2.63	1.61	1.33	246.8	[272]
DD2Y	N-Y	2.76	2.63	1.78	1.04	240.4	[273]
DDLS(30)-Y	N-Y	2.59/2.80	2.49/2.70	1.58	1.19	264.9	[258]
DDLS(50)-Y	N-Y	2.66/2.87	2.58/2.78	1.95	1.19	260.7	[258]
DDLS(70)-Y	N-Y	2.65/2.88	2.58/2.75	1.89	1.19	277.8	[258]
DNS	N-Y	2.41	2.43	1.01	1.51	323.8	[368]
FSU2H*	N-Y	2.76	2.69	1.50	1.37	291.5	[289]
FSU2H*L	N-Y	2.65	2.58	1.90	1.15	287.5	[309]
FSU2H*U	N-Y	2.70	2.64	1.30	1.33	303.3	[309]
QMC-A	N-Y	2.81	2.77	1.04	1.54	195.3	[256]
DD2Y Δ 1.1 – 1.1	N-Y- Δ	3.02	2.88	2.52	1.08	184.5	[251]
DD2Y Δ 1.1 – 1.2	N-Y- Δ	3.25	3.10	3.25	1.15	133.7	[251]
DD2Y Δ 1.2 – 1.3	N-Y- Δ	2.78	2.63	1.79	1.00	240.4	[251]
SFHOY	N-Y	/	3.43	/	/	102.8	[248]
APR	N	3.51	3.494	/	1.65	86.04	[378, 379]
DD2	N	2.68	2.69	/	1.75	261.5	[338, 271]
DD2F	N	3.18	3.17	/	1.56	130.0	[338, 380]
DDLS(30)-N	N	2.60	2.66	/	1.78	268.8	[258]
DDLS(50)-N	N	2.60	2.65	/	1.75	270.3	[258]
DDLS(70)-N	N	2.62	2.60	/	1.73	295.6	[258]
DSH Fiducial	N	3.38	3.40	/	1.62	96.91	[389]
DSH Large Mmax	N	2.78	2.79	/	1.77	181.1	[389]
DSH Large SL	N	3.44	3.42	/	1.60	87.74	[389]
DSH Large R	N	3.08	3.09	/	1.68	139.1	[389]
DSH Small SL	N	3.28	3.27	/	1.65	110.6	[389]
DSH Smaller R	N	3.64	3.64	/	1.57	76.73	[389]
FSU2R	N	2.69	2.70	/	1.76	221.9	[259]
FTNS	N	3.31	3.38	/	1.69	101.6	[333, 381]
GS2	N	2.60	2.58	/	1.78	248.5	[336]
LPB	N	3.14	3.15	/	1.59	130.8	[382, 383]
LS220	N	3.06	3.04	/	1.47	158.4	[156]
LS375	N	2.48	2.46	/	1.58	371.1	[156]
SFHO	N	3.41	3.39	/	1.54	103.4	[271, 370]
SFHX	N	3.11	3.08	/	1.74	134.6	[271, 370]
SRO(Sly4)	N	3.43	3.51	/	1.73	90.90	[384, 385]
TM1	N	2.29	2.32	/	1.75	394.1	[384, 385]
TMA	N	2.50	2.54	/	1.67	297.8	[388, 387]

Table C.4: EoSs used to simulate the asymmetric binary mergers with a total mass of $2.8M_{\odot}$ and a binary mass ratio of $q = 0.8$. First to seventh column give the acronym of the models, hadronic species that the EoS is taking into account, dominant postmerger frequency using a full 3D EoS table, dominant postmerger frequency $f_{\text{peak}}^{1.75}$ using a cold slice supplemented with a constant thermal index $\Gamma_{\text{th}} = 1.75$, ratio ρ_{max}/ρ_0 between the maximum density reached in the remnant 5 ms after the merger and the onset density of heavy baryons, the time-mass averaged thermal $\bar{\Gamma}_{\text{th}}$, and tidal deformability $\Lambda_{1.65}$ of a star with mass of $1.65M_{\odot}$. The last column provides the references for the EoS models. Some of the EoSs produce two dominant peaks with a comparable strength. The table has been published in Ref. [369].

Appendix D

Gravitational wave spectra

In Chapter 4, we presented the results of our BNS merger simulation campaign. It was explained that if the initial masses of the stars in the binary does not exceed the threshold mass, a meta-stable hyper-massive NS is formed. In those simulations, the GW spectra, which is obtained as an output of the simulations, have a form like the one presented in Fig. D.1. The figure shows the effective GW amplitude along the polar direction at a distance of 20 Mpc from the source. Two distinct spectra are shown that are obtained from the two different simulation runs that we have performed with each EoS; the 3D and the $\Gamma_{\text{th}} = 1.75$ run.

The different spectral features are labeled in the figure: the dominant frequency peak f_{peak} , and the secondary features f_{2-0} and f_{spiral} . For the 3D spectrum, the parabolic fit around the three points with highest amplitude is also shown. The position of the maximum of the parabolic fit curve represents the value of f_{peak} . The same procedure is followed for determining the value $f_{\text{peak}}^{1.75}$, as well as for all secondary spectral features.

The part of the graph that corresponds to low frequencies ($f < 1$ kHz) is dominated by waves emitted during the inspiral phase. Knowing that the simulations start only a few orbits before merging, it is clear that the actual intensity emitted in this frequency regime is severely underestimated.

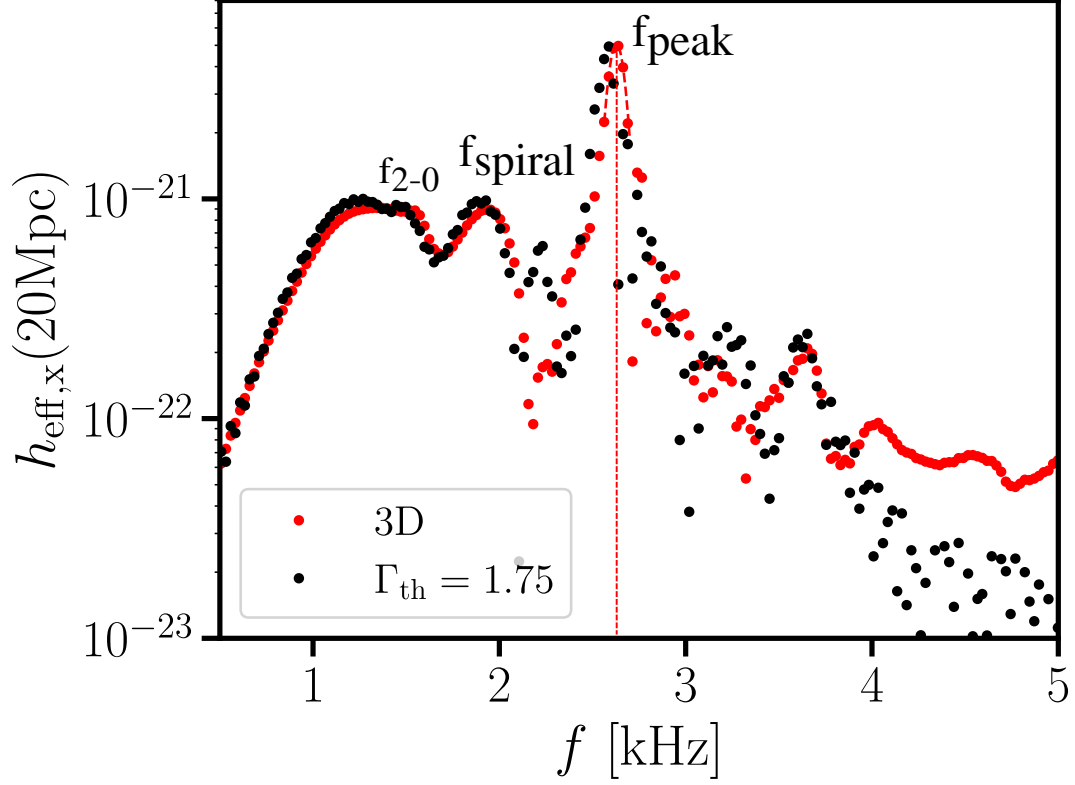


Figure D.1: Typical GW spectrum of the cross polarization shown for the two runs used in the simulation campaign presented in the thesis. Red dots correspond to the spectra obtained from the 3D run, while the black dots are obtained from the $\Gamma_{\text{th}} = 1.75$ run. The three most characteristic peaks are labeled. For the 3D run the parabolic fit of the three points with highest amplitude is also shown. The vertical red line corresponds to the place of the obtained value of f_{peak} .

Appendix E

Toy model of nuclear and hypernuclear matter with low thermal index

In Chapter 4 we investigated in detail the thermal properties of hypernuclear matter with many different EoSs. We then compared the results with the ones obtained using nucleonic models. We found that the characteristic drop of the thermal pressure in hypernuclear matter was directly related with the production of thermal hyperons. That property translated into characteristic features of the BNS merger observables when hypernuclear matter is presented like the increased value of the dominant frequency peak or the enhanced mass ejecta.

However, in some "ab-initio" many-body studies that employ modern nuclear interactions derived from chiral effective field theory (χ EFT) [406, 407, 408], a drop of the thermal pressure is also observed even when only nuclear matter is presented. The authors related the drop of the thermal pressure with the increase of the nucleons effective mass in matter. The behavior of the nucleonic χ EFT EoSs is especially interesting since if the drop of the thermal pressure and thermal index is a generic feature for a wide range of densities, charge fractions and temperatures, it can mimic the effects that we have attributed to the appearance of hyperons. However, it is well known that the χ EFT can not generate an EoS that covers the whole density and temperature regime needed to simulate a BNS merger event. Instead, the approach is applicable only up to densities of around $2\rho_0$ and for temperatures $T \lesssim 30$ [487]. Hence, those EoSs can not be directly used in relativistic simulations. Furthermore, no exotic baryons have been considered as degrees of freedom of the constructed EoSs [406, 407, 408] so comparison between the properties of hypernuclear and nuclear matter can not be done.

To test what happens with the thermal properties of hypernuclear matter when the nuclear matter itself experiences a drop in the thermal pressure, we developed a toy model. The

toy model is based on the RMF framework, and it is built by extending the DDME2 model of Ref. [270] to include hyperons, as it was done in Ref. [254]. We specifically modified the nucleonic part of the DDME2 model to make the thermal index of nuclear matter decrease with density, similar to the χ EFT predictions and the phenomenological models made to replicate that behavior [488]. We achieve the desired behavior of the EoS by changing the functional density dependence of the σ -meson coupling for nucleons. In particular, we used the following density dependence for the σ -meson coupling:

$$g_{\sigma}(\rho_B) = 10.54 \cdot \exp\left(-0.2 \left(\frac{\rho_B}{\rho_0} - 1\right)^2\right). \quad (\text{E.1})$$

With this change, we forced the model to predict a nucleonic effective mass behavior that has a minimum close to ρ_0 , which, as a consequence, significantly reduces the thermal index of matter. The other constants have the same values as in the original DDME2 model. It is important to note, however, that this toy model does not align with nuclear physics and astrophysical constraints and it only serves our purpose of isolating and studying the thermal effects of hyperons.

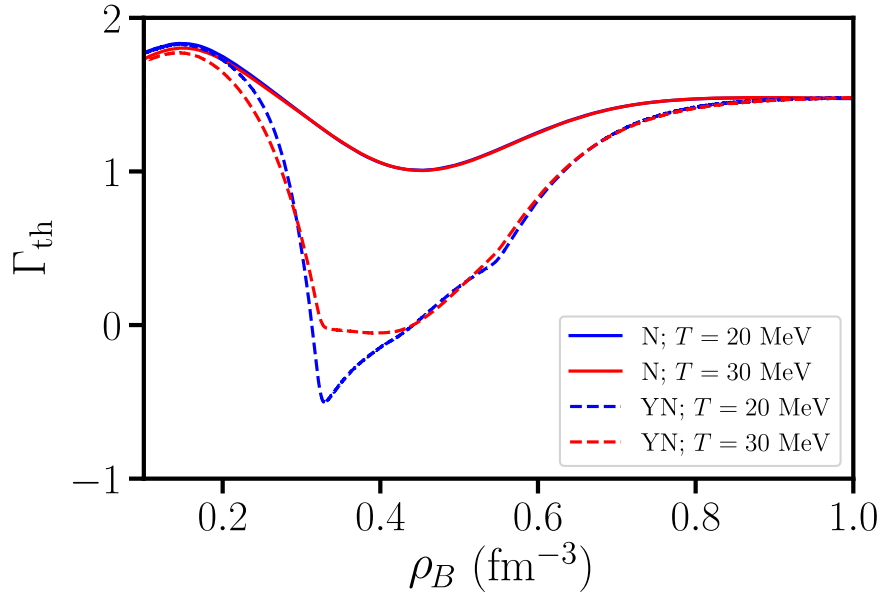


Figure E.1: Γ_{th} for nucleonic (solid lines) and hypernuclear matter (dashed lines) at two different temperatures, $T = 20$ MeV (blue lines) and $T = 30$ MeV (red lines). The figure was published in [364].

With the developed model, two different sets of calculations are done. The first set (labeled with "YN"), calculates the thermal index Γ_{th} of beta-stable hypernuclear matter

at two different constant temperatures ($T = 20$ MeV and $T = 30$ MeV). The second set (labeled with "N"), for the same conditions, calculates Γ_{th} for beta stable nuclear matter, by artificially prohibiting the appearance of hyperons in matter. The results of the thermal index Γ_{th} as a function of the density are shown in Fig. E.1. We can see that indeed the nucleonic model, by construction, produces a dip in the thermal index. However, when the hyperons are allowed to appear in matter they induce an additional decrease of Γ_{th} with respect to the nucleonic benchmark model. We conclude that even if the nucleonic EoS has a soft thermal behavior, the appearance of thermal hyperons will still trigger an additional loss of degeneracy pressure, which results into an even softer thermal behavior.

It is important to address some limitations of our current toy model. As previously mentioned, this simplified model does not accurately reflect the properties of nuclear matter or align with astrophysical observations. Therefore, future research should focus on developing a more realistic parametrization that will take into account the constraints of the EoS. Additionally, in the RMF framework used to construct our model, the effective masses of hyperons and nucleons are linked by symmetry relations between their coupling constants. This means that their effective masses share the same functional dependence, which may not be accurate and could affect the size of the "additional hyperonic drop". However, based on our findings, we can safely conclude that even if a nucleonic model predicts a low thermal index, the presence of hyperons should lead to a further decrease. In future, a realistic model should be built to confirm the size of this effect.

List of Figures

1.1	Scheme of a massive star's structure at the end of its life as a star on the main sequence. The layers' thickness is not drawn to scale.	6
1.2	Evolution path of a PNS ($M < 25M_{\odot}$) after a successful explosion. It is assumed that the star will not reach the mass for a black hole collapse at any stage. The figure is inspired by similar figures in [74] and in [75] (as a chapter of [76]).	8
1.3	Scheme of the structure of a massive cold NS. Please note that the exact thickness of each layer depends on the mass of the star and the interaction assumed between the species (see the discussions in Chapter 3 and Chapter 4.)	9
1.4	Baryon octet. With S the strangeness of the baryons is labeled, while Q represents the baryon electric charge.	11
1.5	Scheme of a BNS merging event in which a hyper-massive NS is formed. Three different phases—inspiral, merger, and the GW phase of the process are marked. The strain of the GWs as a function of time is also shown (amplitude and oscillations not to scale).	15
2.1	Composition of baryonic matter at constant charge fraction and constant temperature. The results obtained with the FSU2H*U model are labeled with solid lines, while the dashed lines correspond to the results obtained with the FSU2H*L model. Different species of matter are represented by each color. The figure is published in [289].	31
2.2	Pressure P , energy per particle E/A and entropy S/A , as functions of baryon density for hypernuclear matter at different charged fractions and temperatures. Dashed lines indicate calculations using the FSU2H*L model, and solid lines indicate calculations using the FSU2H*U model. The figure is published in [289].	33

2.3	Thermal energy density and thermal pressure as functions of the baryonic density for different temperatures, expressed in MeV units. The solid lines label the cases where hyperons are not present in the matter, while the dashed ones represent the calculations where hyperons are taken into account. The figure was published in [289].	37
2.4	Thermal index as a function of the baryon density for various temperatures, expressed in MeV units. The calculations performed without hyperons are marked with solid lines, while the ones that do account for the appearance of hyperons are labeled with dashed lines. The horizontal dashed-dotted orange lines correspond to the thermal index of a non-relativistic ($\Gamma = 5/3$) and ultra-relativistic ($\Gamma = 4/3$) diluted ideal gas, respectively. The figure was published in [289].	39
2.5	P_{th} as a function of ϵ_{th} for various temperatures, expressed in MeV units. The calculations without hyperons are marked with solid lines, while the ones with hyperons are labeled with dashed lines. The dotted lines represent two linear approximations of the $P_{\text{th}}(\epsilon_{\text{th}})$ relation ($\Gamma = 2$ -upper line and $\Gamma = 1.5$ -lower line). The figure was published in [289].	40
2.6	$\Gamma = \Gamma(\rho_B, T)$ fit as a function of the baryonic density ρ_B for the neutrino-free case, without (left panel) and with hyperons (right panel). Dots are representing the exact calculations, while the solid lines are the results from the fits. The temperatures shown are in MeV units. The figure was published in [289]	41
3.1	Temperature (left plot) and electron fraction profiles (right plot) of the stars used in our calculations. Solid lines represent the results with the FSU2H* model, dashed lines represent the ones obtained with the FSU2H*L model, and dashed-dotted lines are results with the FSU2H*U model. Please read the text for detailed motivation of the shown profiles. For all three models, the simulation-inspired profiles at $t = 1$ s are identical. The figure was published in [309].	45
3.2	$M(R)$ curves of cold, β -equilibrated NSs, calculated using the FSU2H* family models. The observational constraints obtained from $2M_{\odot}$ observations [113, 114, 115, 116, 327] and NICER measurements of radii within a 1σ error band [127, 128, 129, 130] are also shown. The figure was published in [309].	47
3.3	$M(R)$ relations of hot stars with temperature and density profiles as described in the text. Results from our three models are shown: FSU2H* model (solid lines), FSU2H*L (dashed lines), and FSU2H*U (dashed-dotted lines). The figure was published in [309].	48

- 3.4 $\Lambda(M)$ relation computed with our three parametrizations together with the observational constraint on the tidal deformability using the universal relations (EoS insensitive) and the spectral EoS obtained from the GW170817 event [351, 352], presented in Ref. [350]. The hatched regions represent the values at 68% confidence level, while the filled boundaries give the values at 95% confidence level. The figure was published in [309]. 51
- 3.5 $\Lambda(M)$ relations of hot stars with temperature and lepton fraction as described in the text. Solid lines represent the results with the FSU2H* model, dashed lines represent the FSU2H*L model, and dashed-dotted lines are results with the FSU2H*U model. The figure was published in [309]. 52
- 3.6 Moment of inertia as a function of the star's mass, obtained with our three parametrizations FSU2H*, FSU2H*L and FSU2H*U. The observational constraints obtained from Ref. [357] are also shown. The figure was published in [309]. 53
- 3.7 Dimensionless moment of inertia obtained with our three parametrizations FSU2H* (solid lines), FSU2H*L (dashed lines), and FSU2H*U (dashed-dotted lines), for PNSs with temperature and lepton fraction as described in the text. The figure was published in [309]. 54
- 3.8 Normalized strangeness number s as a function of M/R , computed with our three parametrizations. The dot at the end of each curve indicates the maximum mass star configuration. The graph was published in [309]. 55
- 3.9 Normalized strangeness number s , computed with our three parametrizations: FSU2H* (solid lines), FSU2H*L (dashed lines), and FSU2H*U (dashed-dotted lines), for PNSs with temperature and lepton profiles as described in the text. The dot at the end of each curve indicates the maximum mass star configuration. The graph was published in [309]. 56
- 4.1 Mass-radius relations computed with many different hyperonic (green lines) and nucleonic (black lines) models. The graph was published in [364], and it was made by Sebastian Blacker. 59
- 4.2 Onset density of heavy baryons as a function of temperature, predicted by the different models utilized in the study. The onset density is defined as the density at which the abundance of exotic species in matter is larger than 1%. The figure was published in [369]. 63
- 4.3 Thermal energy density ε_{th} as a function of density at fixed charge fraction $Y_Q = 0.1$ and temperature $T = 25$ MeV. Three different nucleonic models (solid line for SFHO, dashed line for FSU2R, and dash-dotted line for DD2) are also shown. The figure was published in [369]. 65

4.4	Thermal pressure P_{th} as a function of density at fixed charge fraction $Y_Q = 0.1$ and temperature $T = 25$ MeV. Three different nucleonic models (solid line for SFHO, dashed line for FSU2R, and dash-dotted line for DD2) are also shown. The figure was published in [369].	66
4.5	Thermal index Γ_{th} as a function of density at fixed charge fraction $Y_Q = 0.1$ and temperature $T = 25$ MeV. Three different nucleonic models (solid line for SFHO, dashed line for FSU2R, and dash-dotted line for DD2) are also shown. The figure was published in [369].	67
4.6	Hyperon excess $\Delta\rho_Y$ as a function of density at fixed charge fraction $Y_Q = 0.1$ and temperature $T = 25$ MeV. The figure was published in [369].	68
4.7	Specific heat at constant pressure c_p as a function of density at temperature $T = 25$ MeV and fixed charge fraction $Y_Q = 0.1$. Three different nucleonic models (solid line for SFHO, dashed line for FSU2R, and dash-dotted line for DD2) are also shown. The figure was published in [369].	70
4.8	Time evolution of the maximum rest mass density in units of nuclear saturation density, $\rho_{\text{max}}(t)/\rho_0$ for different hyperonic EoSs. Three different nucleonic models (solid line for SFHO, dashed line for FSU2R, and dash-dotted line for DD2) are also shown. The figure was published in [369].	73
4.9	Time evolution of the fraction of hyperons in the remnant for different hyperonic EoSs. Please note the logarithmic scale on the y-axis. The figure was published in [369].	74
4.10	Time evolution of the average thermal index for different hyperonic EoSs. Three different nucleonic models (solid line for SFHO, dashed line for FSU2R, and dash-dotted line for DD2) are also shown. The value of $\Gamma_{\text{av}}^{\text{th}} = 1.75$ is marked with a horizontal blue line. The figure was published in [369].	75
4.11	Time evolution of the average temperature for different hyperonic EoSs. Three different nucleonic models (solid line for SFHO, dashed line for FSU2R, and dash-dotted line for DD2) are also shown. The figure was published in [369].	76
4.12	Flowchart of the strategy implemented in the simulation campaign. The difference between an observable X obtained from the 3D run and the $\Gamma_{\text{th}} = 1.75$ run will be labeled as ΔX	79

- 4.13 Relative frequency shift as a function of the peak frequency for all EoSs from both sets. Purely nucleonic models are indicated by plus signs, models incorporating hyperons are marked with crosses, and models that include both hyperons and Δ baryons are represented by asterisks. The color bar indicates the ratio between the maximum rest-mass density achieved in the merger remnant within the first 5 milliseconds after the merger and the zero-temperature onset density of heavy baryons ($\rho^{\max}/\rho_{\text{onset}}$). The figure was published in [369]. 80
- 4.14 Relative frequency shift as a function of the mass- and time-averaged thermal index in the remnant for all EoSs from both sets. The notation is the same as in Fig. 4.13. The figure was published in [369]. 81
- 4.15 Relative frequency shift as a function of the mass- and time-averaged hyperon fraction $\bar{Y}_{\text{hyp}}$ in the remnant. Models that incorporate hyperons are marked with crosses, while models that include both hyperons and Δ baryons are represented by asterisks. The color bar indicates the ratio between the maximum rest-mass density achieved in the merger remnant within the first 5 milliseconds after the merger and the zero-temperature onset density of heavy baryons ($\rho^{\max}/\rho_{\text{onset}}$). The figure was published in [369]. 82
- 4.16 Average and maximum deviation for purely nucleonic EoSs from quadratic least-squares fits $f_{\text{peak}} - \Lambda_M$. The solid line shows the behavior of the maximum deviation as a function of the reference mass, while the dashed line shows the functional dependence of the average deviation. The point M at which the maximum deviation reaches a minimum is labeled with a red vertical line. The figure was published in [369]. 85
- 4.17 $f_{\text{peak}} - \Lambda_{1.65}$ relation for all EoSs considered in this work. The notation of the symbols and the color bar is the same as in Fig. 4.13. The solid black line represents the quadratic fit obtained from the purely nucleonic EoS data. The gray band represents the maximum residual of these nucleonic models with respect to the fit. The figure was published in [369]. 86
- 4.18 $f_{\text{peak}} - \rho_{\text{max}}$ relation for all EoSs considered in this work. The notation is the same as in Fig. 4.13. The figure was published in [369]. 87
- 4.19 f_{peak} as a function of stellar parameters of cold nonrotating NS: a) $R_{1.4}$, b) $R_{1.8}$, c) $\Lambda_{1.4}$ and d) $\Lambda_{1.8}$, for the chosen subset of nucleonic and hyperonic models that share the same nucleonic interaction. Purely nucleonic models are indicated by plus signs, models incorporating hyperons are marked with crosses, and models that include both hyperons and Δ baryons are represented by asterisks. The color bar shows the fraction of hyperons in the remnant. The figure was published in [369]. 89

4.20	$\Delta f_{2-0}/f_{2-0}$ as a function of f_{peak} (left panel) and $\bar{\Gamma}_{\text{th}}$ (right panel). The notation is the same as in Fig. 4.13. The figure was published in [369].	93
4.21	$\Delta f_{\text{spiral}}/f_{\text{spiral}}$ as a function of f_{peak} (left panel) and $\bar{\Gamma}_{\text{th}}$ (right panel). The notation is the same as in Fig. 4.13. The figure was published in [369].	94
4.22	Functional dependence between f_{2-0} (left panel) and f_{spiral} (right panel) with f_{peak} . The figure has been published in [369]	95
4.23	Dynamical mass ejecta M_{ej} as a function of the radius $R_{1.4}$ of $1.4 M_{\odot}$ cold NS, 10 ms after merging. Purely nucleonic models are indicated by plus signs, models incorporating hyperons are marked with crosses, and models that include both hyperons and Δ baryons are represented by asterisks. The coloring of the data points signals if the heavy baryons represent more than 1% of the baryon population (green points) or not (blue points). Nucleonic models are labeled with black pluses. The black solid line represents a second-order polynomial fit to the purely nucleonic models. The figure was published in [369].	96
4.24	Relative difference between M_{ej} and a parabolic fit $M_{\text{ej}}^{\text{fit}}(R_{1.4})$ to the purely nucleonic models. The figure was published in [309].	97
4.25	Relative difference between the mass ejecta obtained in the 3D run and the one obtained in the $\Gamma_{\text{th}} = 1.75$ run. The notation is the same as in Fig. 4.23.	98
4.26	Threshold binary mass M_{thres} as a function of M_{max} and $R_{1.6}$. The faded gray plane shows the bilinear fit as described in the text. The small lines suggest the distance of each point to the plane. Purely nucleonic models are indicated by plus signs, models incorporating hyperons are marked with crosses, and models that include both hyperons and Δ baryons are represented by asterisks. The figure was published in [369].	100
4.27	2D plot of M_{thres} as a function of $aM_{\text{max}} + bR_{1.6}$. The coefficients in the linear term $aM_{\text{max}} + bR_{1.6}$ are obtained using the bilinear fit shown in Fig. 4.26. The straight line is the projection of the plane. The figure was published in [369].	101
4.28	Same as Fig. 4.27, but for the $\Gamma_{\text{th}} = 1.75$ run. The figure was published in [369].	102
4.29	ΔM_{thres} as function of M_{thres} . Purely nucleonic models are indicated by plus signs, models incorporating hyperons are marked with crosses, and models that include both hyperons and Δ baryons are represented by asterisks. If two or more models produce equal ΔM_{thres} and M_{thres} , the corresponding data point is labeled with a thicker symbol. The figure was published in [369].	103

- 4.30 Δf_{peak} as a function of f_{peak} in the case of more massive binaries (see the text for detail). For each EoS model, the gray line connects the results of the canonical mass simulations with the higher mass simulations. The arrow indicates the direction towards larger masses. The color bar indicates the ratio between the maximum rest-mass density achieved in the merger remnant within the first 5 milliseconds after the merger and the zero-temperature onset density of heavy baryons ($\rho^{\text{max}}/\rho_{\text{onset}}$) for the more massive configuration. The figure was published in [369]. . . . 105
- 4.31 Δf_{peak} as a function of $\bar{\Gamma}_{\text{th}}$ in the case of more massive binaries (see the text for detail). The notation is the same as in Fig. 4.30. The figure is published in [369]. 106
- 4.32 $\Delta f_{\text{peak}}/f_{\text{peak}}$ as a function of f_{peak} obtained from the asymmetric campaign. Purely nucleonic models are indicated by plus signs, models incorporating hyperons are marked with crosses, and models that include both hyperons and Δ baryons are represented by asterisks. Three of our models have two dominant peaks present in their GW spectrum, which are roughly equally strong. For those EoSs, both peaks are displayed and connected with a dashed line. The figure was published in [369]. 107
- 4.33 $\Delta f/f_{\text{peak}}$ as a function of $\bar{\Gamma}_{\text{th}}$ obtained from the asymmetric campaign. The notation is the same as in Fig. 4.32. The figure was published in [369]. 108
- 4.34 f_{peak} as function of $\Lambda_{1.65}$ in asymmetric mergers with a binary mass ratio $M_1/M_2 = 0.8$. Purely nucleonic models are indicated by plus signs, models incorporating hyperons are marked with crosses, and models that include both hyperons and Δ baryons are represented by asterisks. The black curve represents the least-squares quadratic fit to purely nucleonic models. The gray band labels the maximum residual of purely nucleonic models from the fit. Three of our models have two dominant peaks in their GW spectrum, which are roughly equally strong. For those EoSs, both peaks are displayed and connected with a dashed line. The figure was published in [369]. 109
- B.1 Composition of cold β -stable matter computed with the additional 18 parametrizations (shaded regions) and with our two extreme models (dashed lines - FSU2H*L; dashed-dotted lines - FSU2H*U). The figure is published in [289]. 119
- B.2 Composition of β -stable matter at $S/A = 2$ computed with the additional 18 parametrizations (shaded regions) and with our two extreme models (dashed lines - FSU2H*L; dashed-dotted lines - FSU2H*U). The figure is published in [289]. 120

B.3	Pressure of cold β -stable matter computed with the additional 18 parametrizations (dashed lines) and with our two extreme models (solid blue line - FSU2H*L and solid red line - FSU2H*U). The figure is published in [289].	121
B.4	Pressure of β -stable matter at $S/A = 2$ computed with the additional 18 parametrizations (dashed lines) and with our two extreme models (solid blue line - FSU2H*L and solid red line - FSU2H*U). The figure is published in [289].	122
B.5	$M(R)$ relation of cold β -stable stars computed with the additional 18 parametrizations (dashed lines) and with our two extreme models (solid blue line - FSU2H*L and solid red line - FSU2H*U). The figure is published in [289].	123
B.6	$M(R)$ relation of hot β -stable stars at $S/A = 2$ computed with the additional 18 parametrizations (dashed lines) and with our two extreme models (solid blue line - FSU2H*L and solid red line - FSU2H*U). The figure is published in [289].	123
B.7	$\Lambda(M)$ relation of cold β -stable stars computed with the additional 18 parametrizations (dashed lines) and with our two extreme models (solid blue line - FSU2H*L and solid red line - FSU2H*U). The figure is published in [289].	124
B.8	$\Lambda(M)$ relation of hot β -stable stars at $S/A = 2$ computed with the additional 18 parametrizations (dashed lines) and with our two extreme models (solid blue line - FSU2H*L and solid red line - FSU2H*U). The figure is published in [289].	124
D.1	Typical GW spectrum of the cross polarization shown for the two runs used in the simulation campaign presented in the thesis. Red dots correspond to the spectra obtained from the 3D run, while the black dots are obtained from the $\Gamma_{\text{th}} = 1.75$ run. The three most characteristic peaks are labeled. For the 3D run the parabolic fit of the three points with highest amplitude is also shown. The vertical red line corresponds to the place of the obtained value of f_{peak}	131
E.1	Γ_{th} for nucleonic (solid lines) and hyperonic matter (dashed lines) at two different temperatures, $T = 20$ MeV (blue lines) and $T = 30$ MeV (red lines). The figure was published in [364].	133

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