

Impact of Magnetic Field Topology on Electromagnetic and Gravitational Waves from Binary Neutron Star Merger Remnants

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We perform general relativistic magnetohydrodynamic (GRMHD) simulations of binary neutron star (BNS) mergers with four distinct magnetic field topologies: (i) a dipole pulsar-like configuration, (ii) a mixed linear superposition of poloidal and toroidal components inside the star, and (iii-iv) two topologies featuring a smooth transition from a confined mixed core to a pulsar-like structure at radii $0.95 R_{\text{NS}}$ and $0.5 R_{\text{NS}}$, with R_{NS} the radius of the star. The latter topologies are explored in BNS merger studies for the first time. We evolve systems with two equations of state (EoS), SLy and WFF1, with ADM masses 2.7 and 2.6, respectively, and include an additional lower-mass SLy binary to probe the behavior of long-lived remnants. We perform an extensive analysis of the emission properties of the systems, both electromagnetic and gravitational waves, and of the properties of the remnants, namely their frequency modes, density eigenfunctions, rotation, temperature, and convective stability. We report three key results: (1) for the first time, we assess the convective stability of magnetized remnants, extending previous unmagnetized analyses; (2) we identify a clear secondary peak in the gravitational-wave spectrum of pulsar-like configurations, consistent with the nonlinear coupling of the $m = 0$ and $m = 2$ modes, which is absent in other topologies; and (3) the magnetic field topology strongly influences the gravitational wave emission properties to the extent that nearby (< 50 Mpc) events could allow one to observationally distinguish between different field structures with future gravitational-wave detectors. Across all models, we obtain luminosities compatible with short gamma-ray bursts (sGRBs), with purely poloidal configurations being the most efficient in driving possible relativistic jets.

I. INTRODUCTION

The coincident detection of gravitational waves (GWs) from GW170817, and their electromagnetic (EM) counterparts across the spectrum, marked the beginning of multimessenger astronomy. This event, along with GRB170817A, confirmed that binary neutron star (BNS) mergers can be the progenitors of the central engine that powers short gamma-ray bursts (sGRBs) [1–6]. Notice that, although the progenitor of GW170817 has been officially classified as a BNS, since the estimated individual masses are consistent with the observed mass distribution of neutron stars (NSs) and black holes (BHs) in binaries [7–9], observations cannot rule out the possibility of a neutron star-black hole (NSBH) system. However, these values fall below the minimum BH mass limit determined from X-ray binary observations and stellar collapse models [10]. While mechanisms for the formation of BHs with smaller masses have been proposed [11], observational evidence for BHs below $3M_{\odot}$ is sparse, supporting the classification of GW170817 as a BNS merger.

This single multimessenger event has been used to impose tight constraints on the equation of state (EoS) at supranuclear densities [12–16]. It provided evidence of ejecta masses of approximately $0.01M_{\odot} - 0.05M_{\odot}$ with velocities of $0.1c - 0.3c$, consistent with the production of elements via the rapid neutron-capture process (r-process) needed to explain Milky

Way abundances [17–19]. Additionally, GW170817 offered an independent measure of the Hubble constant [20]. This event underscored the importance of comparing multimessenger observations with theoretical models to understand extreme matter physics, highlighting the necessity of numerical modeling due to the complexity of the underlying phenomena.

The first numerical studies of BNSs began in the late 20th century, primarily focusing on understanding the basic dynamics of BNS mergers, while largely excluding the complexities introduced by magnetic fields or any microphysical processes [21–24]. Early simulations included some degrees of relativistic gravity and simple hydrodynamic models [25–27]. With advancements in computational capabilities and numerical methods, the 2000s saw the introduction of more sophisticated simulations incorporating full general relativity (GR) [28, 29] and magnetohydrodynamics (MHD) [30]. These improvements allowed the study of increasingly complex initial conditions [31–37], more realistic EoS [38], and higher numerical resolution [39, 40]. Despite the progress made in BNS merger simulations, there remains a need to enhance the modeling of magnetic reconnection, integrate more detailed microphysical processes, and deepen our understanding of the role of magnetic instabilities, and in particular, the role of the magnetic turbulence in the stability properties of the BNS remnant. Understanding these instabilities is crucial for interpreting the expected multimessenger signals from BNS systems, and may offer new insights into the mechanisms behind the launching of relativistic jets. However, magnetic fields significantly increase the complexity of simulations and in-

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introduce new computational challenges. To accurately capture magnetic instabilities, simulations with high numerical resolution are required. This has often rendered studies of magnetic instabilities unfeasible, due to the high computational costs of such simulations (see [41] for a review). To overcome the need for high resolution, alternative approaches, explored in [42–48], have been proposed. They involve using large-eddy simulations and subgrid-scale models. These methods compensate for limited resolution by introducing additional terms in the evolution equations to approximate unresolved subgrid dynamics. A key advantage is that they allow for the use of weaker magnetic fields, closer to those observed in pulsars, and are considered to be more effective in modeling the small-scale, quasi-isotropic turbulent fields generated by the Kelvin-Helmholtz instability (KHI) [44]. These approaches are potentially preferred over artificially strong initial poloidal fields, which, as argued in [47], may produce unrealistic results. However, a notable drawback is that the outcomes may be highly sensitive to the choice of subgrid model and its parameterization, raising concerns about whether these models accurately reflect the true physical processes.

Our GRMHD numerical simulations in [49–51] with numerical resolutions ranging from $\Delta x \sim 90 - 220$ m have shown that these mergers can produce a collimated, mildly relativistic outflow confined by a helical magnetic field emanating from the poles of a BH surrounded by a magnetized accretion disk formed after the collapse of the hypermassive neutron star (HMNS) remnant. This outflow, which is driven by the Blandford-Znajek (BZ) mechanism [52], was identified as an incipient jet. We used irrotational NSs of equal mass and a polytropic EoS with a polytropic index $\Gamma = 2$. Similar results were found in NSBH mergers [53, 54]. Follow-up studies included variations such as initial star spins [55], different magnetic dipole orientations [56], realistic piecewise polytropic EoSs (SLy and H4) [57], and an M1 neutrino transport scheme [58]. Ciolfi et al. [59, 60] conducted BNS merger simulations with numerical resolutions $\Delta x \sim 177 - 220$ m resulting in a long-lived supramassive neutron star (SMNS) remnant, finding a collimated outflow but with heavy baryon pollution that prevents the formation of an outflow compatible with sGRBs. By contrast, Mösta et al. [61] employed numerical resolutions $\Delta x \sim 55 - 220$ m to perform pure hydrodynamic simulations of BNSs. The post-merger remnant was endowed with a poloidal magnetic field and reported the formation of a mildly relativistic jet, with neutrino cooling reducing baryon pollution and increasing the Lorentz factor, though still below sGRB levels. Most recently, Kiuchi et al. [62] performed very high-resolution simulations with $\Delta x \sim 12$ m and noted the formation of a mildly relativistic jet, with potential for a higher Lorentz factor if efficient Poynting flux conversion occurs, despite severe baryon loading. Aguilera-Miret et al. [63] using a resolution of $\Delta x \sim 120$ m with a subgrid model observed a helical magnetic field but found no other evidence of jet launching.

Most previous numerical studies assumed that the stars are initially endowed with dipole magnetic field, either confined entirely within the star’s interior or extending from the interior to the exterior, similar to the standard picture for pulsars’

exterior magnetic fields.

However, the interior magnetic field topology of NSs remains unknown. Observations of pulsar J0030+0451 with NASA’s Neutron star Interior Composition Explorer (NICER) have revealed unexpected hotspot locations, with up to three hotspots detected in the southern hemisphere, instead of the expected one at each pole [64–66]. These observations suggest the presence of a more complex magnetic field topology. Moreover, both Newtonian and GRMHD simulations have shown that purely toroidal and purely poloidal magnetic field configurations are unstable [67–75], consistent with analytical predictions [76–78].

NICER’s observations and the stability requirements for NS magnetic fields strongly advocate for the adoption of more sophisticated models to characterize their topology. For example, it has been suggested that NSs are born with both poloidal and toroidal fields exceeding 10^{14} G, but the poloidal component decays by a factor of 10 over 10^6 yr [79]. As a first step, one can consider a superposition of toroidal and poloidal magnetic fields as the most general configuration.

Multiple equilibrium solutions comprising mixed-field configurations have been proposed [80]. Among these, the twisted-torus is a promising candidate for modeling magnetic fields in NS interiors [81]. In this configuration, the poloidal component of the field extends to the exterior of the star, while the toroidal one is confined within a torus inside the star, ensuring the continuity and smoothness of field lines at the stellar surface. However, numerical simulations of isolated self-consistent rotating NSs with various twisted-torus configurations showed that stability depends strongly on the mixture of the different EM field topologies, with some configurations being more unstable than the others [74]. Ongoing efforts have been devoted to the development of realistic NS magnetic field models [82–88], with recent studies looking at the impact of these topologies on jet launching, the emitted gravitational wave signal, the global amplification of the magnetic field, and the nature of the ejected material [63, 89–91].

Driven by the anticipated multimessenger observations of BNSs, we probe here the impact of distinct magnetic field topologies on the stability properties of the remnant, on the jet launching, and more generally on both the EM and GW signals. We consider four different topologies: (i) a dipole poloidal pulsar-like configuration, (ii) a mixed linear superposition of poloidal and toroidal components inside the star, and (iii–iv) two topologies with a smooth transition from a confined mixed core to a pulsar-like structure at radii $0.95R_{\text{NS}}$ and $0.5R_{\text{NS}}$ (with R_{NS} the equatorial radius of the neutron star), ensuring physical plausibility and continuity across the stellar interior and magnetosphere. All configurations have exclusively poloidal fields outside the star, recovering the typical pulsar picture. Together, these span a broad range of plausible interior topologies. We find that systems with pulsar-like magnetic fields consistently exhibit the largest magnetic energies. These cases are consistent with observed sGRBs and kilonovae phenomena, expected to originate as the aftermath of BNS mergers [18]. In scenarios where the remnant is a NS whose lifetime exceeds our simulated time, the magnetic field topology modifies the temperature distribution, regions

of convective stability, and the remnant’s structure. Consistent with [89, 92], we find that the f_2 GW peak from magnetized BNSs is shifted compared to purely hydrodynamic cases, with the shift being highly dependent on the magnetic topology. We do not consider neutrino effects here, as our focus is on magnetic field topology. However, a more complete analysis should include them, as neutrinos may drive jet formation without magnetic fields [93] and reduce baryon pollution (see e.g. [58, 94]).

The remainder of the paper is organized as follows. In Sec. II, we summarize the numerical methods to solve the GRMHD equations, the initial data used to model the BNSs, and the grid structure, along with a suite of diagnostics used to verify the reliability of our numerical calculations. Our main results are reported in Sec. III. Finally, we summarize our results and conclude in Sec. IV. Additional comments on the magnetic power spectra and convective stability of the BNS remnants are provided in Appendices A and B, respectively. Throughout the paper, we adopt geometrized units ($G = c = 1$) except where otherwise explicitly stated. Greek indices denote all four spacetime dimensions, while Latin indices imply spatial components only.

II. METHODS

The numerical methods employed to evolve our BNS models are the same as those in [95, 96]. In this section, we briefly introduce our notation and summarize our numerical schemes, along with the initial data. For more detailed information, we refer the reader to [95, 96] and the references therein.

A. Formulation and numerical scheme

To solve the Einstein equations coupled with the equations of ideal MHD in a dynamic, curved spacetime, we utilize the Illinois GRMHD code. This code is the basis for its publicly available counterpart embedded in the Einstein Toolkit [97, 98] and is integrated into the Cactus computational framework [99]. The Illinois GRMHD code employs Carpet for moving mesh refinement [100, 101], and it has been extensively tested and applied in various GRMHD contexts, including isolated BHs with magnetized accretion disks and BNS mergers undergoing delayed collapse (for a review, see [95, 96]). It solves the Baumgarte–Shapiro–Shibata–Nakamura (BSSN) formulation [102, 103] to evolve the spacetime metric. This is coupled with the puncture gauge conditions, specifically the “1+log” slicing for the lapse function and the hyperbolic gamma-driver condition for the shift cast in first order form (see Eq. (2)–(4) in [104]), using fourth order centered spatial differencing, except on shift advection terms, where a fourth order upwind differencing is used. To mitigate high-frequency modes that may arise during the evolution, we apply fifth-order Kreiss-Oliger dissipation to all evolved variables, with a dissipation strength of 0.18.

The ideal MHD equations are solved in a flux conservative formulation, via high-resolution shock-capturing (HRSC) schemes (see Eqs. (27)–(29) in [97]), that employs the piecewise parabolic method (PPM) [105] coupled to the Harten, Lax, and van Leer (HLL) approximate Riemann solver [106]. As a standard procedure to avoid the breakdown of these schemes we impose a tenuous atmosphere with rest-mass density $\rho_{\text{atm}} = 10^{-10} \rho_{0,\text{max}}$, where $\rho_{0,\text{max}}$ is the initial maximum value of the rest-mass density of the system.

Instead of evolving the magnetic field, we evolve the vector potential A^μ by integrating the magnetic induction equation which ensures that the “no monopoles” constraint is satisfied. We adopt the generalized Lorenz gauge, described in [107], to avoid the build-up of spurious magnetic fields. Finally, we impose “outflow” boundary conditions to the primitive variables and the vector potential [104]. The time integration is performed with the method of lines using a fourth-order accurate Runge-Kutta method, with a Courant-Friedrichs-Lewy (CFL) factor of 0.45.

B. Initial data

We consider BNS configurations in quasi-equilibrium circular orbits that merge to form either a stable remnant, below the supramassive limit, or a transient HMNS that will undergo delayed collapse into a BH. The binaries consist of two identical irrotational NSs, modeled by a piecewise polytropic representation of the Skyrme-Lyon (SLy) [108] and the WFF1 [109] nuclear EoSs. The initial data are computed using the Parallel Compact Object CALculator (PCOCAL) code [110, 111], and the properties of the configurations considered are summarized in Table I.

We note that these two EoSs broadly align with the current observational constraints on NSs. For example, the maximum mass configuration of an isolated star predicted by SLy is $M_{\text{sph}}^{\text{max}} = 2.05 M_\odot$, while that predicted by WFF1 is $M_{\text{sph}}^{\text{max}} = 2.12 M_\odot$. Both are consistent with: i) $M_{\text{sph}}^{\text{max}} > 2.072^{+0.067}_{-0.066} M_\odot$ from the NICER and XMM analysis of PSR J0740+6620 [112]; ii) $M_{\text{sph}}^{\text{max}} > 2.01^{+0.017}_{-0.017} M_\odot$ from the NANOGrav analysis of PSR J1614-2230 [113]; iii) $M_{\text{sph}}^{\text{max}} > 2.01^{+0.14}_{-0.14} M_\odot$ from the pulsar timing analysis of PSR 0348+0432 [114]; and iv) $M_{\text{sph}}^{\text{max}} > 2.14^{+0.20}_{-0.18} M_\odot$ from the NANOGrav and the Green Bank Telescope [115]. We also note that SLy predicts that a star with a mass of $1.4 M_\odot$ has a radius of $R = 11.46$ km, consistent with the value $R = 11.94^{+0.76}_{-0.87}$ km obtained by a combined analysis of X-ray and GW measurements of PSR J0740+6620 [116]. Such a star modeled with WFF1 has a radius of $R = 10.23$ km, which is slightly below the constraint imposed by [116]. Consistent with this, the combined analysis of the LIGO/Virgo/KAGRA (LVK) collaboration of the progenitors of GW170817 with the radio-timing observations of the pulsar J0348+0432 [114, 117] constrains the radius of a NS with mass in the range $1.16 M_\odot - 1.6 M_\odot$ to be $11.9^{+1.4}_{-1.4}$ km at the 90% credible level. However, the NICER analysis of PSR J0030+0451 [64] constrains the radius of a NS with mass of $1.44^{+0.15}_{-0.14} M_\odot$ to be

TABLE I. Summary of the initial properties of the BNS configurations. We list the EoS employed to model the NSs, the equatorial coordinate radius R_{NS} toward the companion, the compactness $\mathcal{C}$, the tidal deformability $\Lambda = (2/3)k_2\mathcal{C}^{-5}$ (where k_2 is the second Love number), the rest-mass of the binary M_0 , the ADM mass M , and the angular velocity Ω for an initial binary coordinate separation of ~ 45 km.

EoS	$R_{\text{NS}}[\text{km}]$	$\mathcal{C}$	Λ	$M_0[M_\odot]$	$M[M_\odot]^a$	$M\Omega$
SLy	9.05	0.17	353	3.04	2.70	0.030
SLy	9.25	0.16	498	2.83	2.54	0.029
WFF1	7.86	0.19	232	2.95	2.60	0.037

^a We note that the threshold gravitational mass at infinite separation for prompt collapse is $2.83M_\odot$ and $2.68M_\odot$ for SLy and WFF1, respectively [120].

$R = 13.02^{+1.24}_{-1.06}$ km. The SLy and WFF1 EoSs predict that a $1.44M_\odot$ star has a radius of 11.45 km and 10.24 km, respectively. Therefore, SLy is slightly below the NICER constraint, while WFF1 deviates by at least 14%. Consequently, NICER analysis favors stiffer EoSs. Furthermore, the LVK analysis of GW170817 predicts that the tidal deformability of a $1.4M_\odot$ NS is $\Lambda_{1.4} = 190^{+390}_{-120}$ at the 90% credible level [117]. For such a star, the tidal deformability values are $\Lambda_{1.4} = 306.4$ for SLy and $\Lambda_{1.4} = 157.33$ for WFF1. Hence, LVK analysis favors softer EoSs, such as SLy and WFF1 [118]. On the other hand, the analysis of GW190814 [119] suggests a mass of $2.59^{+0.08}_{-0.09}M_\odot$ for the secondary object, which favors stiffer EoS. Note, however, that the constraints imposed by event GW190814 are more uncertain as there is no confirmation that the secondary object is a neutron star.

C. Grid structure

Details on the grid structure are given in Table II. It consists of two sets of nested refinement boxes, each centered on one of the NSs. Each set contains nine boxes, with sizes and resolutions differing by factors of two. When two boxes overlap, they are merged into a single box centered on the center of mass of the binary system. The finest box surrounding each NS has a half-length on the side of approximately 1.25 times the initial equatorial radius R_{NS} of the NS. This configuration allows us to resolve the NS equatorial radius with $\gtrsim 90$ grid points similar to our previous studies in [51, 89]. Since we do not expect reflection symmetry across the equatorial plane to strongly impact the outcome of our simulations [56], we impose it to reduce computational costs.

TABLE II. Grid structure for models listed in Table I. The computational mesh consists of two sets of 9 nested refinement boxes. The finest boxes are centered on each star and have a half-length of $\sim 1.25R_{\text{NS}}$, where R_{NS} is the initial equatorial coordinate radius toward the companion. The number of grid points covering the equatorial radius of the NS is denoted by N_{NS} . In all cases, we impose reflection (equatorial) symmetry about the orbital plane. For the lower-mass SLy case, we include the ADM mass in the EoS name to distinguish it from the higher-mass case.

EoS	Grid hierarchy [†]	Max. resolution	N_{NS}
SLy	3024.03 km/ 2^{n-1}	95 m	91
SLy2.54	5748 km/ 2^{n-1}	90 m	102
WFF1	2646.03 km/ 2^{n-1}	86 m	91

[†] The spatial extent of the simulation box for each refinement level, where n denotes the level number.

D. Magnetic field models

To mitigate the effects of magnetic winding during the inspiral phase, we evolve the configurations in Table I with no magnetic fields to approximately two orbits ($\lesssim 3$ ms, roughly smaller than the Alfvén timescale; see below) before merger, except for the lower-mass SLy configuration.¹ At that point, the NSs are threaded by a dynamically weak magnetic field generated by the vector potential

$$\mathbf{A} = (1 - a_T) \mathbf{A}_P + a_T \mathbf{A}_T, \quad (1)$$

where $\mathbf{A}_P$ generates a poloidal magnetic field component given by

$$\mathbf{A}_P = \frac{\pi \varpi^2 I_P r_P^2}{(r_P^2 + r^2)^{3/2}} \left[1 + \frac{15 r_P^2 (r_P^2 + \varpi^2)}{8 (r_P^2 + r^2)^2} \right] \hat{\mathbf{e}}_\phi, \quad (2)$$

that approximately corresponds to that generated by a current loop with radius r_P inside the star, where I_P is the current, $r^2 = \varpi^2 + z^2$, with $\varpi^2 = (x - x_{\text{NS}})^2 + (y - y_{\text{NS}})^2$ and $(x_{\text{NS}}, y_{\text{NS}}, 0)$ is the coordinate location of the center of mass of the NS. The vector potential $\mathbf{A}_T$ generates a tilted toroidal magnetic field component and is given by

$$\mathbf{A}_T = 2 I_T \log \left(\frac{2 r_T}{\varpi^2} \right) \max(P - P_{\text{cut}}, 0)^{n_b} \hat{\mathbf{e}}_y, \quad (3)$$

where r_T and I_T are the loop radius and the current that generates the toroidal field, analogous to r_P and I_P , respectively. The cutoff pressure parameter P_{cut} confines the toroidal component to reside within $P > P_{\text{cut}}$. Note that

¹ Note that the lower-mass SLy configuration (which we denote by SLy2.54) initially threaded by a pulsar-like poloidal magnetic field (see Eq. (2)), and reported previously in [51], is considered in our study to probe the effect of magnetic fields in a long-lived remnant.

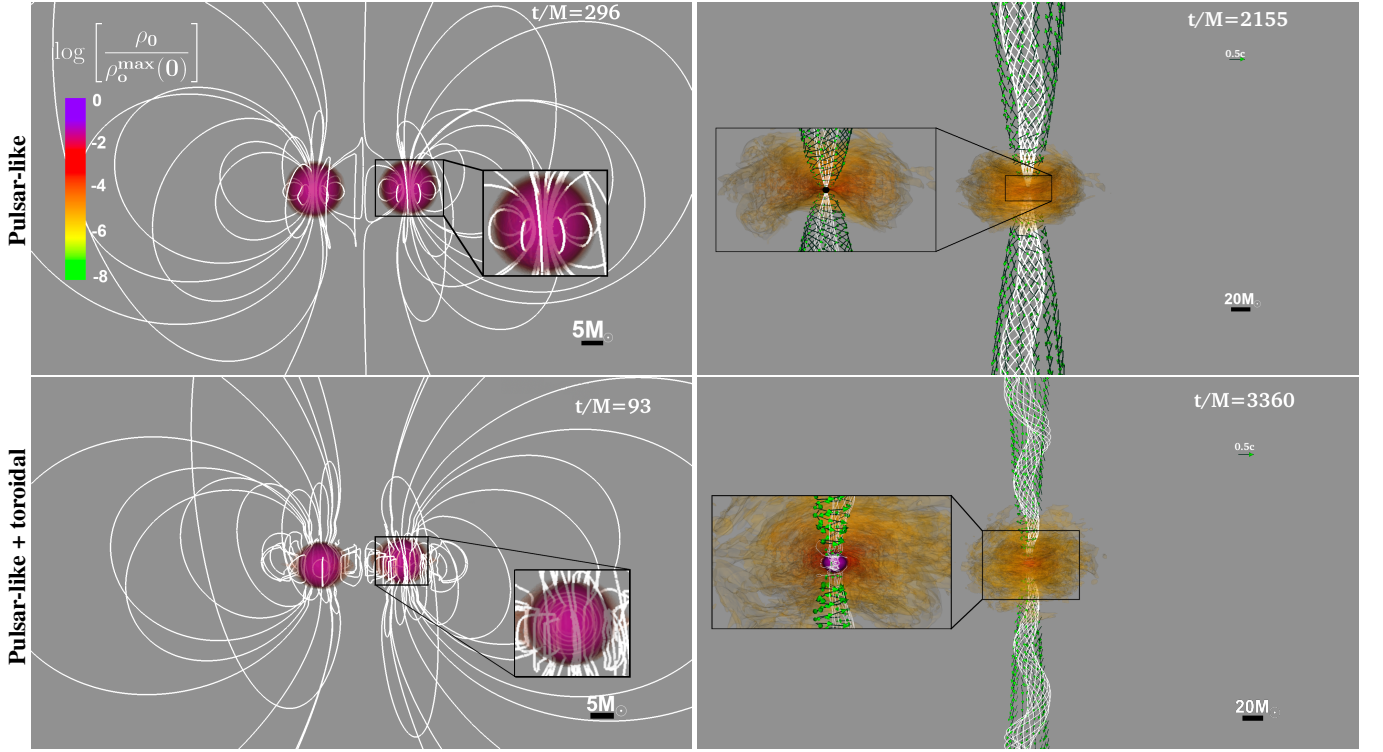


FIG. 1. Volume renderings of the rest-mass density, normalized to its initial maximum value (logarithmic scale), at the moment of magnetic field insertion (left) and after the BNS remnant reaches quasiequilibrium (right). Top panel shows the SLy BNS with a pure poloidal (pulsar-like) magnetic field, while the bottom panel depicts the WFF1 BNS with a mixed magnetic field inside the star ($r_1 = 0.95 R_{\text{NS}}$, see Sec. IID), smoothly transitioning to an external dipole field. Insets highlight the initial magnetic field configuration inside the star (left), and the helical magnetic structure after an incipient jet is launched from the pole of the binary’s remnant (right top panel BH remnant and right bottom NS remnant). White lines represent magnetic field lines, arrows indicate plasma velocities, and the black sphere marks the BH apparent horizon. Here, $M = 2.7M_{\odot}$ for the top panels and $M = 2.6M_{\odot}$ for the bottom panels, with $M_{\odot} = 4.9 \times 10^{-3} \text{ ms} = 1.4 \text{ km}$.

in contrast to [89], the toroidal component is chosen to be perpendicular to the orbital angular momentum, rather than aligned with it, hence being called “tilted”. This means that the vector potential generating this component lies within the orbital plane. This component generates a mixed magnetic field, with both poloidal and toroidal components. The parameter n_b determines the degree of central condensation of the magnetic field.

We consider the following magnetic field topologies:

- **Pulsar-like magnetic field (P):** Following [96], the stars are endowed with a purely poloidal magnetic field, characterized by $a_T = 0$ in Eq. (1) (see top-left panel in Fig. 1). This serves as our fiducial case.
- **Superposition of poloidal + toroidal magnetic field (SP):** The stars are endowed with a mixed magnetic field configuration within their interior, consisting of a linear superposition of toroidal and poloidal components, and in the exterior we impose a pulsar-like magnetic field. In this case, we set $a_T = 0.5$ in Eq. (1).
- **Tilted toroidal magnetic field in the NS + external pulsar-like field (T):** The stars are endowed with a mixed magnetic field within their interior, which

smoothly transitions to a pulsar-like magnetic field extending from the interior to the exterior (see bottom-left panel in Fig. 1). To achieve a smooth transition between the interior and exterior magnetic field components, we set $a_T = e^{-(r/r_1)^{2p}}$ as in [121], where r_1 and p are free parameters that determine the degree of central condensation of the interior magnetic field component, with r_1 indicating where the transition occurs, and p setting how fast this transition is. For this magnetic field configuration, we consider two cases: i) $r_1 = 0.95 R_{\text{NS}}$, and denote it $\mathbf{T}_{0.95}$, which corresponds to an essentially mixed field confined in the star; and ii) $r_1 = 0.5 R_{\text{NS}}$, where the mixed field is more confined to the star’s core, which is denoted as $\mathbf{T}_{0.5}$. In both cases, we set $p = 10$.

The only difference between configurations **SP** and **T** is the value of a_T , which in case **T** ensures a smooth transition and also allows to confine the toroidal component (in the **SP** configuration it is confined by the pressure, guaranteeing that there is no toroidal component outside the stars), whereas the confinement in configuration **T** is determined by r_1 . The equation for the vector potential is given in Eqs. (1), (2), and (3), with a_T the one stated for each configuration.

To differentiate between the different cases, we labeled them as EoS_magnetic-field-configuration. For the confined

TABLE III. Initial magnetic field energy $\mathcal{M}$ in erg and maximum magnetic field strength $B_{\max}$ in G for all BNS configurations in Table I in which all cases have a maximum value of $\beta^{-1} \equiv P_{\text{mag}}/P = 0.003125$ within the star, except for the SLy2.54_P configuration in which $\beta^{-1} = 0.0023$.

Case	$\mathcal{M}[\text{erg}]$	$B_{\max}[\text{G}]$
SLy_P	$10^{50.66}$	$10^{16.23}$
SLy2.54_P	$10^{50.61}$	$10^{16.20}$
SLy_SP	$10^{48.72}$	$10^{15.26}$
SLy_T0.95	$10^{50.53}$	$10^{16.17}$
SLy_T0.5	$10^{49.96}$	$10^{15.88}$
WFF1_P	$10^{50.82}$	$10^{16.41}$
WFF1_SP	$10^{48.86}$	$10^{15.43}$
WFF1_T0.95	$10^{50.39}$	$10^{16.19}$
WFF1_T0.5	$10^{50.05}$	$10^{16.02}$

mixed cases, an additional tag is used to indicate the confinement level of the tilted toroidal component (see Table III). Additionally, for the lower-mass SLy binary, we include a tag for its ADM mass, to distinguish it from the heavier configurations with the same EoS. For example, SLy_P denotes a BNS configuration modeled with the SLy EoS, ADM mass of $2.7M_{\odot}$, and a purely poloidal magnetic field, while SLy_T0.5 represents the same binary but with a mixed magnetic field confined within the star, with $r_1 = 0.5 R_{\text{NS}}$ (see Sec. IID). Cases with no magnetic fields (hydrodynamical only) are denoted by the tag H.

To ensure a fair comparison of the evolution across the different magnetic field configurations, we choose the above free parameters such that the maximum magnetic-to-gas pressure ratio within the interior of the NSs is $\beta^{-1} \equiv P_{\text{mag}}/P = 0.003125$ across all cases, as in [49, 50], which corresponds to a maximum magnetic field as measured by a normal observer of $B_{\max} \lesssim 10^{16.4}$ G (see Table III). This value is significantly higher than the surface magnetic field strengths of $\sim 10^8 - 10^{12}$ G typically observed in NSs within binary systems, as inferred from binary pulsar observations [122, 123]. Following [49, 50], we choose this value to mimic the amplified field strengths expected from the exponential growth of magnetic instabilities, which has been observed in high resolution BNS simulations [45, 124] (see Sec. IIE for a more detailed discussion on magnetic instabilities). We find that the magnetic power spectra of our BNS remnants are in broad agreement with higher-resolution and large-eddy simulations [40, 43] (see Appendix A).

We also consider a slightly weaker pulsar-like magnetic field configuration by setting $\beta^{-1} \equiv P_{\text{mag}}/P = 0.0023$ for the SLy2.54_P case, as in [51]. As shown below, the remnant in this case behaves similarly to those modeled with the WFF1 EoS (it does not collapse), allowing us to analyze the stability properties of longer-lived NS remnants.

E. Magnetic instabilities

The KHI acts on small scales through the small-scale turbulent dynamo mechanism [125, 126]. The turbulent flow resulting from the shearing of fluid flowing in opposite directions stretches and folds the magnetic field lines. The stretching increases the length of the field lines, enhancing the magnetic field tension along the direction of stretching due to the frozen-in property of magnetic fields in a high-conductivity fluid. In highly turbulent regions, magnetic diffusion acts to dissipate the magnetic energy. Eventually, the magnetic field grows to a point where its back-reaction on the fluid flow becomes significant. This back-reaction modifies the turbulence, reducing the efficiency of the dynamo, reaching saturation. The growth rate of the KHI, σ_{KHI} , is inversely proportional to the wavelength [125, 126]

$$\sigma_{\text{KHI}} \sim \frac{\Delta v}{\lambda}, \quad (4)$$

where Δv is the difference in velocity across the fluid layers that come into contact and move past each other. Eq. (4) suggests that the small-scale vortices grow faster than the large-scale vortices. However, in numerical simulations the wavelength is limited by the numerical resolution. Note that an initial magnetic field strength of 10^{11} G requires a numerical resolution of at least ~ 7 m to properly capture the magnetic field amplification due to the KHI [62]. Given the limited computational resources available, we employ stronger initial fields to mimic the magnetic field amplification due to the KHI. During the first 2 ms following merger we observe that, consistent with the values reported in [51], the growth rate of the total magnetic energy is $\sim 1.18 \text{ ms}^{-1}$ for SLy binaries and $\sim 1.37 \text{ ms}^{-1}$ for those modeled with WFF1. The slightly higher amplification factor in the latter cases is likely due to a higher speed difference across the shear surface. As in [51, 62], the measured growth rate is two orders of magnitude lower than the expected $\sigma_{\text{KHI}} \sim 10^2 \text{ ms}^{-1}$ from perturbation theory.

After the BNS remnant settles into a massive central core wrapped by a low-density cloud of matter, the high velocity gradients that drive the KH instability weaken. Amplified magnetic fields then regulate the flow through angular momentum transport, suppressing the conditions necessary for the KHI to continue. The magnetorotational instability (MRI) and magnetic winding emerge during the post-merger phase of BNS mergers, enhancing magnetic turbulence and angular momentum transport [42, 46, 90, 91].

The remnant becomes unstable to the MRI when the condition $\partial\Omega^2/\partial\ln\varpi < 0$, with Ω the angular velocity, is satisfied [127, 128]. Two fluid elements in the remnant, connected by a magnetic field line, experience different rotational speeds. The inner element, closer to the stellar center, rotates faster than the outer one. This differential rotation stretches the magnetic field line, generating magnetic tension. The faster-rotating inner element pulls on the magnetic field line, exerting a force on the slower-rotating outer element, thereby increasing its angular velocity. Conversely, the outer element exerts a retarding force on the inner element, decreasing its an-

gular velocity. This exchange of angular momentum between the elements destabilizes the rotational equilibrium, causing the inner element to move inwards (losing angular momentum) and the outer element to move outwards (gaining angular momentum). This process of stretching the magnetic field lines and exchanging angular momentum intensifies, leading to an increase in magnetic tension and further amplification of the velocity perturbations. Consequently, this feedback loop results in the exponential growth of both the magnetic field and velocity perturbations, rapidly amplifying the strength of the magnetic field. The growth rate and the fastest-growing wavelength of the MRI read

$$\sigma_{\text{MRI}} = \frac{1}{2} \frac{\partial \Omega}{\partial \ln \varpi}, \quad (5)$$

$$\lambda_{\text{MRI}} \approx \frac{2\pi v_A}{\Omega} \approx \frac{2\pi \sqrt{b^P b_P / (b^2 + \rho_0 h)}}{\Omega} \quad (6)$$

$$\sim 18 \text{ cm} \left(\frac{B}{10^{11} \text{ G}} \right) \left(\frac{\rho_0}{10^{14} \text{ g cm}^{-3}} \right)^{-1/2} \left(\frac{\Omega}{10^3 \text{ rad s}^{-1}} \right)^{-1},$$

where v_A is the Alfvén speed, $|b^P| = \sqrt{b^\mu b_\mu - (b_\mu (e_\phi)^\mu)^2}$ and $(e_\phi)^\mu$ is the toroidal orthonormal vector comoving with the fluid.

To determine if the MRI is adequately resolved in our simulations, two key conditions must be satisfied: i) the wavelength of the fastest-growing mode, λ_{MRI} , has to be resolved by at least $\simeq 10$ grid points; and ii) λ_{MRI} has to fit in the remnant [129, 130]. To check condition i), we compute the λ_{MRI} -quality factor $Q_{\text{MRI}} \equiv \lambda_{\text{MRI}}/dx$, with dx the local grid spacing, which measures the number of grid points per fastest growing MRI mode (see top panels of Fig. 2). We observe that the $Q \gtrsim 10$ criterion is satisfied in the binary remnant cases **P** and **T** over a substantial portion of the remnant. However, in the **SP** case the quality factor is lower. Note that fixing the maximum value of the magnetic-to-gas pressure ratio in the stellar interior to $\beta^{-1} = 0.003125$ induces different magnetic field strength depending on the topology of the magnetic field (see Table III). The **SP** configuration presents the lowest magnetic field strength, which helps explain the small values of Q for that case, which is $\lesssim 5$ in the central core of the remnant, and around $\lesssim 10$ in the low density cloud of matter wrapping it. This is due to the fact that the MRI wavelength is proportional to $\rho_0^{-1/2}$ and therefore decreases in regions of higher density. To check condition ii), we plot the rest-mass density normalized to its maximum value along λ_{MRI} on the meridional plane. In all cases, the fastest-growing mode λ_{MRI} fits within the remnant.

Magnetic winding occurs in differentially rotating objects when a poloidal magnetic field (with both radial and vertical components) generates a linear growth of the toroidal magnetic field in a timescale of $\tau_{\text{wind}} \sim \Delta\Omega^{-1}$, where $\Delta\Omega$ is the difference in angular velocity across the star [131]. As the magnetic field lines are wound up and the toroidal component strengthens, magnetic tension develops, resisting differential rotation through magnetic braking. This process eventually drives the system toward uniform rotation on the Alfvén

timescale [132]

$$\tau_A \sim 10^2 \text{ ms} \left(\frac{\rho_0}{10^{15} \text{ g cm}^{-3}} \right)^{1/2} \left(\frac{B}{10^{15} \text{ G}} \right)^{-1} \left(\frac{R}{10^6 \text{ cm}} \right). \quad (7)$$

Note that this timescale is comparable to the lifetime of typical BNS remnants and may influence their subsequent evolution [41, 47]. As emphasized in [41], higher resolution studies than those performed here are necessary to assess the effects of initially weak magnetic fields.

F. Jets consistent with sGRBs

Jet-like structures appear in a variety of natural phenomena across different scales in the universe, and are characterized as highly collimated streams of matter and energy ejected from a source [133]. In the context of stellar compact binary mergers, they are thought to lead to sGRBs, although the central engine is not yet completely understood. There is a general consensus about the possibility of such transients emerging from a BH surrounded by a magnetized accretion disk; however, there is no clear picture of whether sGRBs can emerge from NS binary remnants [41, 96, 134]. However, an increasing number of studies have been dedicated to exploring scenarios where GW170817-like events may leave behind a rapidly rotating NS [51, 62, 135].

Numerical simulations have also shown jet-like structures emerging from stable NS remnants [41, 45, 51, 60, 61], with recent studies focusing on how the α - Ω -dynamo mechanism can lead to large-scale magnetic fields in HMNSs, resulting in a Poynting flux-dominated relativistic outflow similar to that of the BZ mechanism for BHs [41, 136]. The formation of such jets from NSs hinges on dynamo processes influenced by microscopic parameters such as magnetic helicity.² Some studies suggest that finite helicity can suppress dynamo actions, hindering the organization of magnetic fields into large-scale structures [138–140]. Additionally, inefficiencies in outflow acceleration for neutron star remnants have been observed in GRMHD simulations [41, 51, 60, 61]. So, the question remains, are these incipient jets consistent with sGRBs?

To gain a better understanding of what systems can be the progenitors of the central engine that powers sGRBs, it is useful to study which jet properties are consistent with these EM signals. sGRBs are thought to originate from the internal and external shocks in a highly relativistic jet, with Lorentz factors of at least $\Gamma_L \gtrsim 20$, with typical values of $\Gamma_L \sim \mathcal{O}(10^2)$, and 90% of the total γ -ray count is observed in $\lesssim 2 \text{ s}$ [141–145]. The terminal Lorentz factor Γ_∞ can be estimated near the base of the outflow by (see [51] for a more detailed discussion)

$$\Gamma_\infty \approx \frac{\mu}{2} \approx \frac{\sigma}{2} \approx \frac{b^2}{2\rho_0} = \frac{\rho_B}{\rho_0}, \quad (8)$$

where μ is the ratio between the energy flux and the rest-mass energy flux, σ is the magnetization parameter, and

² A new formalism for the gauge-invariant computation of magnetic helicity transport in numerical relativity codes has been presented in [137].

$\rho_B = b^2/2 = B_{\text{co}}^2/8\pi$ (where $b^\mu := B_{\text{co}}^\mu/\sqrt{4\pi}$ for comoving magnetic field B_{co}^μ) is the EM energy density. Whereas BH central engines demonstrate efficient acceleration capabilities (see e.g. [49, 94, 146]), the same is not as clear for stable NS remnants. Thus, understanding whether NS jets can emit sGRBs requires a rigorous assessment of baryon pollution, which is not within the scope of this work.

Despite lasting less than 2 s, sGRBs release an extraordinary amount of energy, with typical isotropic-equivalent luminosities ranging from $\sim 10^{49} - 10^{52} \text{ erg s}^{-1}$ [133, 147, 148]. Therefore, computing the isotropic-equivalent luminosities provides another method to assess whether the jets we observe in our simulations are capable of producing sGRBs. In addition, the duration of these transients correlates with the lifetime of the accretion disk (the jet's fuel) surrounding the remnant, which can be estimated by computing the accretion rate (see for instance [49]).

G. Diagnostics

Throughout the entire numerical evolution of our configurations, we monitor various quantities to ensure the reliability of our simulations and to track the evolution of the BNS remnants. A brief description of these quantities is provided below.

a. Constraints: Throughout the evolution, we monitor the L_2 norm of the normalized Hamiltonian and the normalized momentum constraints as defined in Eqs. (40)-(43) of [104]. In all simulations the Hamiltonian constraint violations remain smaller than $\sim 0.05\%$ during the inspiral. In the case of BH formation, the constraints peak at $\sim 1.6\%$ and then relax to $\sim 0.1\%$ as the system reaches a steady-state. The momentum constraint violations remain smaller than $\sim 1\%$ during the inspiral, peak at $\sim 10.5\%$ during BH formation (if applicable) and then relax to $\sim 0.8\%$ as the system reaches a steady-state. These values are similar to those previously reported in our long-term, hydrodynamic simulations of spinning BNSs modeled using the SLy and ALF2 EoSs [149].

b. BH Horizons: In cases where the binary undergoes delayed collapse, we employ the `AHFinderDirect` thorn [150] to track the apparent horizon of the BH. Additionally, we estimate the BH mass and its dimensionless spin using the isolated horizon formalism as described in [151].

c. Gravitational waves: To compute the Weyl scalar Ψ_4 , we use the `Psikadelia` thorn, decomposing it into $s = -2$ spin-weighted spherical harmonic modes. Using Eqs. (2.8), (2.11), and (2.13) in [152], we calculate the total flux of energy and angular momentum carried away by GWs at ten extraction radii, ranging from $r_{\text{min}} \approx 120M$ and $r_{\text{max}} \approx 840M$ and perform an extrapolation to infinity. We find that between $\sim 1.5\%$ and $\sim 3.0\%$ of the total energy of our binaries is radiated away in the form of gravitational radiation, while between $\sim 20.0\%$ and $\sim 27.9\%$ of the angular momentum is radiated (see Table IV).

Following [153, 154], we analyze the post-merger gravitational wave signals using a modified version of the Prony's method, which allows us to accurately characterize the main

gravitational wave mode and their frequencies. We use the `ESPRIT` (Estimation of Signal Parameters via Rotational Invariance Techniques) variant of Prony's method [155, 156], which is robust against noise and involves fitting a larger number of complex exponentials to the data. Our methodology uses a moving window interval of 4 ms to perform the analysis over time. This process includes performing Fast Fourier Transforms (FFT) on segments of the signal, extracting Fourier amplitudes at selected fixed frequencies, and correlating these with specific oscillation modes. We report, in Table IV, the fundamental frequency associated with the dominant $l = m = 2$ oscillation mode f_2 , its transient value f_{2i} , typically computed during the first few milliseconds after merger, along with the f_{max} , which corresponds to the instantaneous GW frequency at merger, i.e. at the GW peak amplitude (see the vertical dashed lines in Fig. 9) [157]. Furthermore, we obtain the time-frequency spectrograms of the $l = m = 2$ component of the GW strain for all models, and plot them against the active modes that are responsible for the most part of the GW emission. We also performed a density distribution FFT and extracted the Fourier amplitudes at specific fixed frequencies corresponding to the main modes observed in the spectrograms. These amplitudes represent the eigenfunctions of the oscillation modes [154].

We use time-domain waveforms from the binaries in Table I that do not collapse to a BH, injecting them into Gaussian noise consistent with the sensitivity of the third-generation Einstein Telescope (ET). The signals are then reconstructed using `BayesWave` [158], a Bayesian data analysis algorithm that employs a model-agnostic approach for signal recovery.

d. Global Conservation: We check the conservation of the ADM mass and the ADM angular momentum along with the conservation of the rest-mass given by the expressions (21)-(23) in [95], along with the conservation of the rest-mass computed through Eq. (23) in [95]. The ADM mass is conserved to $< 1\%$, the angular momentum to within $\lesssim 3\%$, and the rest-mass within $\lesssim 1.2\%$.

e. Electromagnetic signatures: We compute the magnetic energy [95]

$$\mathcal{M} = \int u^\mu u^\nu T_{\mu\nu}^{(\text{EM})} dV, \quad (9)$$

and monitor it throughout the simulations. where u^μ is the fluid 4-velocity and $T_{\mu\nu}^{(\text{EM})} = b^2 u_\mu u_\nu + b^2 g_{\mu\nu}/2 - b_\mu b_\nu$ is the stress-energy tensor associated with the magnetic field, and dV is proper volume element on a $t = \text{constant}$ spatial slice. Here $b^2 = b^\mu b_\mu$ is the magnetic field measured in the fluid's comoving frame [95].

We evaluate the outflow of matter

$$M_{\text{esc}} = \int_{r>r_0} \rho_* \Theta(-u_t - 1) \Theta(v^r) dx^3 + \int_{t'=0}^t \int_{\delta\mathcal{D}} \rho_* \Theta(-u_t - 1) \Theta(v^r) v^i dS_i dt', \quad (10)$$

where $\rho_* = \sqrt{-g} \rho_0 u^t$, with g the determinant of the 4-metric. The surface element on the sphere is denoted by dS_i . Heaviside functions Θ are used to include only the material with

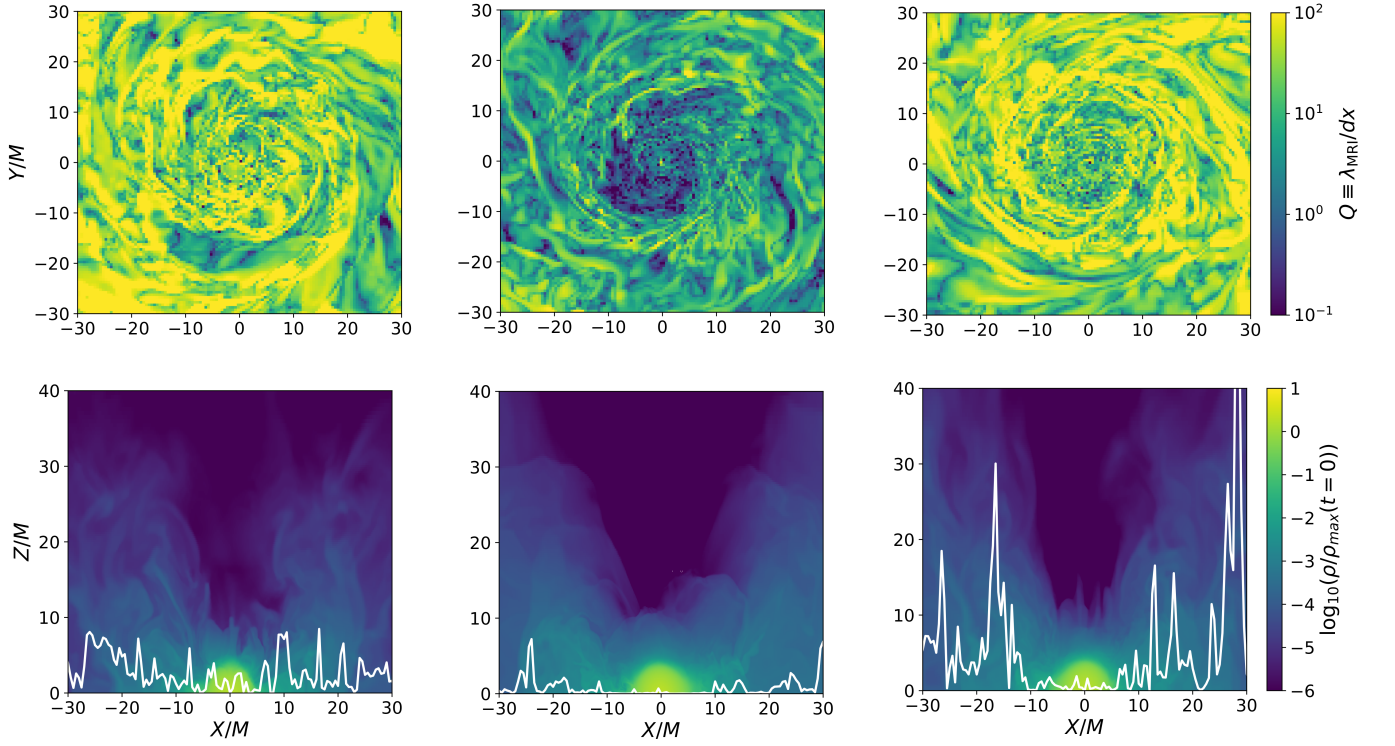


FIG. 2. Top panel: Pseudocolor plot of the quality factor $Q_{\text{MRI}} = \lambda_{\text{MRI}}/dx$ on the equatorial (xy) plane for the remnant of the SLy BNS at ~ 3 ms after merger with: i) a pure initial poloidal field (left); a pulsar-like + mixed magnetic field (middle); and an interior mixed field transitioning smoothly to an external pulsar-like magnetic field at $0.95R_{\text{NS}}$ (right). Bottom panel: Rest-mass density normalized to its initial maximum value (logarithmic scale) on a meridional (xz) slice, along with the λ_{MRI} (white line) for the same cases as in the top panels. For the most part the λ_{MRI} fits within the remnant.

positive specific energy $E = -u_t - 1$ and positive radial velocity v^r . The second term is the contribution from the rest mass leaving the boundary of the simulation domain $\delta\mathcal{D}$. We measured the outflow of matter at different radii between $r_0 = 30M$ and $100M$ and find that the difference between them is less than $\lesssim 3.6\%$ consistent with the values found in BNS simulations modeled with a SLy EoS in [51].

In addition, we compute the fluid luminosity, defined as

$$L_{\text{fluid}} = \int \sqrt{-g}(-T_t^{i(\text{fluid})} - \rho_0 u^i) dS_i, \quad (11)$$

where $T_t^{i(\text{fluid})}$ is the projection of the stress-energy tensor of a perfect fluid, and the EM Poynting luminosity as

$$L_{\text{EM}} = - \int \sqrt{-g} T_t^{i(\text{EM})} dS_i, \quad (12)$$

over ten spherical surfaces with radii ranging from 120 to $840M$. Furthermore, for the SLy $M = 2.70M_\odot$ cases which form black holes, we compute the expected electromagnetic luminosity for a jet powered by the steady-state BZ mechanism, which can be estimated as [159]

$$L_{\text{BZ}} \sim 10^{52} \left(\frac{\chi_{\text{BH}}}{0.64} \right)^2 \left(\frac{M_{\text{BH}}}{2.5M_\odot} \right)^2 \left(\frac{B_{\text{pol}}}{10^{16} \text{ G}} \right)^2 \text{ erg s}^{-1}, \quad (13)$$

where χ_{BH} and M_{BH} are the dimensionless spin and mass of the BH remnant, respectively, and B_{pol} is the strength of the magnetic field at the BH pole.

We compute the effective Shakura–Sunyaev parameter associated with the effective viscosity due to magnetic stresses

$$\alpha_{\text{SS}} \sim \frac{\text{magnetic stress}}{\text{pressure}} = \frac{T_{\hat{r}\hat{\phi}}^{\text{EM}}}{P}, \quad (14)$$

where $T_{\hat{r}\hat{\phi}}^{\text{EM}} = e_{\hat{r}}^\mu e_{\hat{\phi}}^\nu T_{\mu\nu}^{\text{EM}}$ is the component of the electromagnetic stress-energy tensor in the local comoving frame and $e_{\hat{i}}^\mu$ the corresponding basis of local tetrads (see Eq. (26) in [160]). We find that $\alpha_{\text{SS}} \sim 10^{-3}$ inside the remnant, and $\sim 10^{-2}$ in the disk. However, note that it has been shown this value depends on resolution [124, 161, 162].

f. Remnant properties: We estimate the coordinate angular velocity of the metastable HMNS as

$$\Omega := \frac{u^\phi}{u^t}, \quad (15)$$

with u^ϕ and u^t the azimuthal and time components of the four-velocity.

Following [163], we compute the temperature T of the HMNS via

$$\epsilon_{\text{th}} = \frac{3k_B T}{2m_n} + f \frac{aT^4}{\rho_0}, \quad (16)$$

where $\epsilon_{\text{th}} = \epsilon - \epsilon_{\text{cold}}$ is the specific thermal energy density, while ϵ and ϵ_{cold} are the specific energy density and its cold part. Here, k_B is Boltzmann's constant, m_n is the mass of the neutron, and a is the radiation constant. The first term approximates the thermal energy of the nucleons, while the second term accounts for the thermal energy contributed by relativistic particles. The factor f reflects the number of species of relativistic particles that contribute to the thermal energy, and we assume $f = 11/4$, i.e. we account only for the contribution of electrons and positrons to the thermal radiation. At temperatures exceeding tens of MeV, thermal neutrinos are produced abundantly [164], such that the assumption of $f = 11/4$ is no longer valid. A more consistent treatment would instead rely on tabulated equations of state that incorporate thermal, compositional, and neutrino contributions. However, the simple model using a single value of f suffices as a first approximation.

g. The Brunt-Väisälä frequency: To probe the convective stability of our BNS remnants and understand their internal dynamics, we estimate the Brunt-Väisälä frequency $\mathcal{N}$, which quantifies the oscillation frequency of a vertically displaced fluid element around its equilibrium position [165, 166], via

$$\mathcal{N}^2 \propto \mathcal{B}_i \mathcal{G}^i, \quad (17)$$

where $\mathcal{B}_i$ is the Ledoux discriminant, and $\mathcal{G}^i$ is the gravitational acceleration. Convective stability has previously been studied for non-magnetized BNS remnants [154, 167, 168]. It has also been studied in the linear regime of the MRI (see, e.g. [169]). In the nonlinear regime of the MRI, numerical simulations of shearing boxes have been performed in three dimensions in the context of accretion disks [170]. Here, we use the Brunt-Väisälä frequency as a proxy to study the nonlinear regime in the context of magnetized BNS remnants for the first time. A more detailed description of our approach and the validity of our approximations is provided in Appendix B.

III. RESULTS

The dynamics of the inspiral phase of the BNS cases in Table III is roughly the same. The stars begin in a quasi-circular orbit and gradually lose energy and angular momentum through GW radiation, causing their coordinate separation to shrink. As they approach each other, tidal forces become significant, distorting the stars until they merge, forming a dual-core transient configuration. These cores orbit each other within a hot, nuclear-density envelope and eventually merge, forming a highly differentially rotating central core surrounded by a low-density cloud of matter. Depending on the EoS, the total mass of the system, and the magnetic field configuration, the binary remnant undergoes collapse into a BH or forms a metastable HMNS remnant with a lifetime longer than the duration of our simulations. Consistent with [59, 89], we verify that, depending on the initial poloidal magnetic field strength in the NSs, the BH + disk remnant launches a magnetically-driven outflow compatible with sGRBs. Further studies of the baryon pollution in the

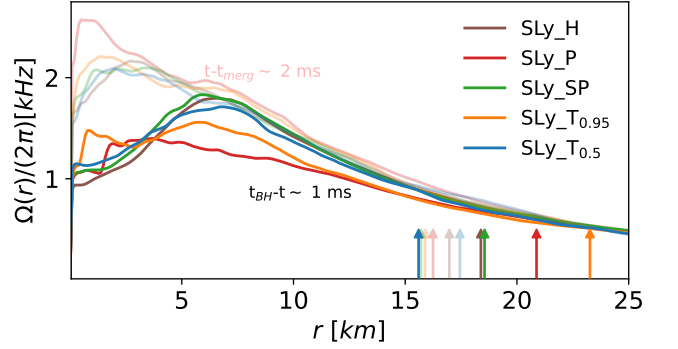


FIG. 3. Average azimuthal angular velocity profiles of the binary remnant that undergo collapse (see Table IV) at ~ 2 ms following merger (faded lines) and at 1 ms before BH formation (bold lines). Arrows marks the rest mass density contour at $\rho_0 = 10^{-3} \rho_{0,\text{max}}(t=0)$, which gives a rough measure of the size of the remnant.

funnel are required to understand if there are other mechanisms for efficient particle acceleration in the stable NS remnants scenario [62]. Key physical properties characterizing the evolution of the binaries are summarized in Table IV.

A. Post-merger evolution

1. Undergoing delayed collapse to a BH

Consistent with previous GRMHD studies using similar magnetic field strengths and resolutions [51, 57], we find that SLy binaries with an ADM mass of $2.7M_\odot$ merge to form a highly differentially rotating remnant with a total rest mass of $\sim 3M_\odot$. Magnetic instabilities, driven mainly by the MRI, induce angular momentum transport from the inner to the external layers of the remnant which eventually causes the formation of a uniformly rotating core surrounded by a Keplerian cloud of low density material. Since the core exceeds the supramassive threshold for the SLy EoS ($2.96M_\odot$) [51, 171], the remnant ultimately collapses to a spinning black hole wrapped by an accretion disk.

Depending on the initial magnetic field configuration (see Table IV), the lifetime of the remnant ranges from ~ 4 to ~ 20 ms. In particular, the **SP** and **T_{0.5}** cases, our least magnetized models making them especially challenging for reliably resolving the MRI (see Table III), have a lifetimes comparable to the non-magnetized scenario. Fig. 2 shows the MRI quality factor $Q_{\text{MRI}} = \lambda_{\text{MRI}}/dx$ on the equatorial plane at $t - t_{\text{merg}} \sim 3$ ms for the **P**, **SP**, and **T_{0.95}** configurations. The **P** and **T_{0.95}** remnants are clearly unstable to MRI. Angular momentum redistribution can be seen in Fig. 3, which shows that on an Alfvén time scale the binary remnant of these cases consists of a roughly uniformly rotating core immersed in a low-density Keplerian envelope.

The total magnetic energy, including contributions from EM outflows leaving the numerical domain, is displayed in Fig. 4. Following merger, the magnetic energy increases by

TABLE IV. Summary of the key properties of the binary mergers. Here, t_{merg} is the merger time in ms, f_{2i} is the instantaneous frequency at the time of the merger associated with the dominant oscillation mode ($l = m = 2$), and f_2 is the frequency once the binary remnant has settled into a quasi-stable state. All values of frequency are reported in kHz. The f_2 values are not reported for cases that collapse ~ 4 ms after merger, as the frequency increases as the remnant becomes more compact, making it impossible to define a single dominant frequency. The fractions of the total energy M and angular momentum J carried off by GWs are E_{GW}/M and J_{GW}/J , respectively. M_{esc} is the escaping rest-mass (ejecta) calculated using Eq. (10) at 20 ms after merger. L_{EM} is the Poynting luminosity in erg s^{-1} 20 ms after merger. L_{knova} , τ_{peak} , and T_{peak} denote, respectively, the estimated peak EM luminosity in erg s^{-1} , rise time in days, and temperature of the potential kilonova arising from the sub-relativistic ejecta in K (see Eqs. (21), (20), and (22)). These are calculated from the ejecta mass and the averaged ejecta velocity v_{eje} of the outflow computed at ~ 160 km from the remnant. t_{BH} is the black hole formation time in ms since merger. M_{BH} and χ_{BH} are the mass and spin of the black holes at the end of the simulation, calculated using the isolated horizon formalism. A dash symbol indicates “not applicable”. Finally, the last column shows the final time post-merger of the simulations.

Label	t_{merg}	f_{2i}	f_2	f_{max}	E_{GW}/M	J_{GW}/J	M_{esc}	$\langle v_{\text{eje}} \rangle$	L_{EM}	L_{knova}	τ_{peak}	T_{peak}	Fate	t_{BH}	M_{BH}	χ_{BH}	t_{final}
SLy_H	7.04	3.61	3.76	1.97	2.69%	27.9%	0.49%	0.21	-	$10^{40.63}$	3.65	$10^{3.28}$	BH	15.00	2.49	0.67	29.4
SLy_P	7.04	-	-	1.98	1.83%	21.3%	1.13%	0.30	$10^{51.08}$	$10^{40.71}$	4.64	$10^{3.28}$	BH	3.34	2.59	0.74	21.6
SLy_SP	7.04	4.46	3.82	1.96	2.59%	27.2%	0.47%	0.32	$10^{50.29}$	$10^{40.59}$	2.90	$10^{3.34}$	BH	15.41	2.48	0.66	20.6
SLy_T0.95	7.02	-	-	1.96	2.07%	22.9%	0.83%	0.25	$10^{50.87}$	$10^{40.61}$	4.35	$10^{3.31}$	BH	3.70	2.59	0.74	21.3
SLy_T0.5	7.04	3.84	3.93	1.97	2.41%	25.8%	1.15%	0.20	$10^{50.92}$	$10^{40.59}$	5.77	$10^{3.30}$	BH	13.17	2.46	0.64	20.9
WFF1_H	3.26	3.88	4.00	2.32	3.01%	27.9%	0.53%	0.21	-	$10^{40.49}$	3.74	$10^{3.35}$	HMNS	-	-	-	40.2
WFF1_P	3.25	4.04	4.11	2.33	2.49%	24.0%	3.50%	0.22	$10^{51.84}$	$10^{40.79}$	9.34	$10^{3.21}$	HMNS	-	-	-	40.3
WFF1_SP	3.26	3.82	3.83	2.34	2.83%	26.8%	0.52%	0.20	$10^{50.84}$	$10^{40.62}$	3.79	$10^{3.28}$	HMNS	-	-	-	40.1
WFF1_T0.95	3.25	3.72	3.83	2.34	2.47%	24.2%	1.35%	0.20	$10^{51.60}$	$10^{40.62}$	6.15	$10^{3.29}$	HMNS	-	-	-	40.0
WFF1_T0.5	3.26	3.92	4.04	2.33	2.73%	26.0%	1.55%	0.21	$10^{51.55}$	$10^{40.66}$	6.35	$10^{3.27}$	HMNS	-	-	-	40.8
SLy2.54_P	6.77	3.34	3.41	1.92	1.55%	20.0%	2.57%	0.22	$10^{52.80}$	$10^{40.89}$	7.92	$10^{3.17}$	SMNS	-	-	-	50.6

more than an order of magnitude within the first ~ 1.5 ms, consistent with amplification driven by the KHI at the shear interface [51]. The total magnetic energy increases exponentially, $\propto \exp(2\gamma_{\text{KH}}(t - t_{\text{merg}}))$, with γ_{KH} the characteristic growth rate. Consistent with [51], we observe that the relative amplification growth rate is $\gamma_{\text{KH}} \sim 1.38 \text{ ms}^{-1}$. As pointed out by [51], this magnetic energy growth rate $2\gamma_{\text{KH}} \sim 2 \text{ ms}^{-1}$ is several orders of magnitude lower than the growth rate predicted by linear perturbation theory $\sigma_{\text{KH}} \sim 10^2 \text{ ms}^{-1}$ based on $\Delta v \sim 0.1c$ and shear layer width $d \sim 400$ m. This is likely due to lack of resolution, which partially resolves the fastest-growing KH modes, and/or the relatively low efficiency of the instability in converting kinetic shear energy into magnetic energy. Studies of spinning and irrotational BNSs using numerical resolutions of a few meters are needed to estimate the growth rate in the non-linear regime.

As the remnant begins to settle, high velocity gradients, responsible for driving the KH instability, diminish and the MRI and magnetic winding take over causing a change in the growth rate of the magnetic field. The magnetic energy in the **P** and **T**_{0.95} remnants peaks roughly at $\gtrsim 10^{51}$ ergs just before BH formation followed by a sharp drop as highly magnetized material in the core is accreted into the BH. Afterward, the magnetic energy decreases more gradually, continuing until the end of the simulation due to the ongoing infall of surrounding material. In **SP** and **T**_{0.5} cases, our least magnetized models, we observe that after $t - t_{\text{merg}} \sim 4$ ms when the KHI weakens, the magnetic field initially grows exponentially mainly due to the MRI and magnetic winding, then it transitions to a non-linear regime before eventually saturating. Following this saturation, the magnetic energy declines,

which is particularly evident in **SP** after $t - t_{\text{merg}} \sim 5$ ms. This decline, also reported in higher-resolution simulations (see e.g. [47, 62]), has been attributed to the conversion of magnetic energy into kinetic energy during the acceleration of ejecta or to the collapse of an unstable magnetic configuration [172]. We note that the angular velocity profile at ~ 1 ms before BH formation is higher in these two cases compared to the more strongly magnetized ones (see Fig. 4). By $t - t_{\text{merg}} \sim 6$ ms the total magnetic energy gradually begins to rise again due to the influence of magnetic instabilities, although this increase is partially balanced by the conversion of magnetic energy into the kinetic energy of the accelerated gas. Following BH formation, the magnetic energy in **T**_{0.5} drops sharply as its highly magnetized central core is swallowed by the newly formed BH. Afterward, the magnetic energy slowly decreases as infall material is accreted. In contrast, during BH formation the magnetic energy in **SP** exhibits a slight decrease (see inset in Fig. 4), followed by a continuous increase. Although the value of the initial magnetic-to-gas pressure ratio within the NSs is constant across all cases, **SP** has the lowest initial magnetization energy and hence the wavelength of the fastest growing MRI mode in the HMNS epoch is relatively short. Following BH formation, the magnetic energy in this case continues to grow until the termination of our simulation, driven by both the MRI and magnetic winding. Since the MRI characteristic wavelength scales with $\rho_0^{-1/2}$ (see Eq. (6)), it is resolved with the adopted resolution.

As shown in Table IV, the mass M_{BH} of the BH remnant and its dimensionless spin χ_{BH} is roughly the same across different cases, with values of $M_{\text{BH}} \sim 2.5$ and $\chi_{\text{BH}} \sim 0.7$ near the termination of the simulations. This is somehow expected

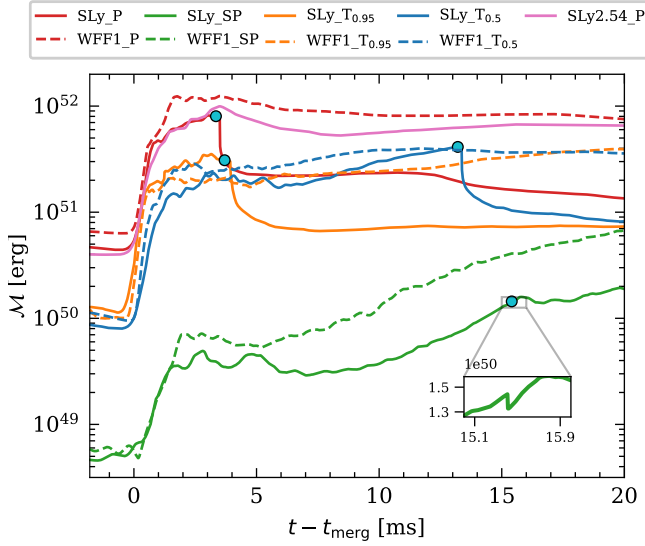


FIG. 4. Evolution of the magnetic energy $\mathcal{M}$ as computed via Eq. (9) for cases listed in Table I endowed with the magnetic field configurations in Sec. IID. The cyan circles mark the BH formation time. The inset highlights a slight decrease in $\mathcal{M}$ just after BH formation in the SLy_SP case. However, its magnetic energy continues to increase beyond this point.

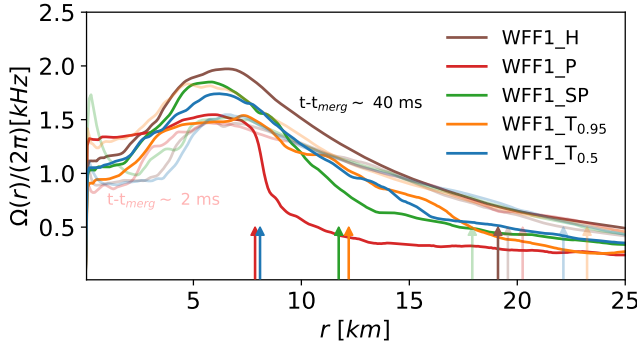


FIG. 5. Average azimuthal angular velocity profiles of the stable binary remnants (see Table IV) in at ~ 2 ms following merger (faded lines) and near to the termination of our simulations (bold lines). Arrows marks the rest mass density contour at $\rho_0 = 10^{-3} \rho_{0,\max}(t = 0)$, which gives a rough measure of the size of the remnant.

as the physical properties of the BH remnant are mainly determined by the mass of the system and the orbital angular momentum of the system at merger.

2. (Meta)stable remnants

For the SLy2.54 and WFF1 binaries, the merger remnants are a stable SMNS [51] and long-lived HMNS remnants, respectively, with lifetime exceeding our simulated time $t - t_{\text{merg}} > 40$ ms. The early evolution of these systems closely resembles that discussed in the previous section, with the key

distinction being the nature of the remnant formed in each case. Following merger, the magnetic energy increases by more than an order of magnitude within the first ~ 1.5 ms. After that point, the magnetic energy remains roughly constant in all cases but WFF1_SP, which is one of the cases with the lowest initial magnetic field energy (recall that we initially fixed the maximum of the magnetic-to-gas pressure ratio inside the stars to be the same for all the magnetic field configurations). In this case, we observe that the magnetic energy continues to grow until the termination of our simulation, likely driven by both the MRI and magnetic winding.

Fig. 5 displays the rotational profiles of the WFF1 remnants as computed via Eq. 15 shortly after merger ($t - t_{\text{merg}} \sim 2$ ms) and near the termination of our simulations ($t - t_{\text{merg}} \sim 40$ ms). Notice how, in the magnetized cases, the curves flatten as time goes on. This flattening of the rotational profile is due to the transport of angular momentum induced by the magnetic viscosity. Naturally, as the employed resolution for the SP case is not enough to fully resolve the MRI, we observe that this case retains a significant degree of differential rotation even at late times. By the same logic, the non-magnetized case is, as expected, the one retaining the largest degree of differential rotation. The pulsar-like configuration is almost rigidly rotating (see solid red line in Fig. 5, where the solid arrow indicates the radius of the corresponding remnant). We observe that the efficiency of the angular momentum redistribution correlates to the maximum initial magnetic field strength (see e.g. Table III), which is expected as the fastest-growing wavelength of the MRI is proportional to the magnetic field strength (see Eq. (6)).

Figs. 6 and 7 show the temperature of the remnants calculated with Eq. (16), on the left of each panel, and the Brunt-Väisälä frequency $\mathcal{N}$ given by Eq. (17) (see Appendix B for more details), on the right. Cyan lines display different rest-mass density contours corresponding to the boundaries between segments of the piecewise equation of state used to model the star, while the green dashed line marks the surface of the star, which we define as $10^{-3} \rho_0^{\max}$.

Immediately after merger, the remnants reach maximum temperatures of $\gtrsim 50$ MeV, due to the strong shock-heating. By 20–40 ms after merger, the cores cool and contract, while the outer layers retain heat in a ring structure, consistent with findings from previous studies of non-magnetized BNS mergers (see, e.g. [173]). Here, we find that the properties of this ring are not only affected by the presence of a magnetic field, but also by its topology.

The hot ring can be attributed to the differential rotation profile present in the remnant shortly after merger. During this epoch, the outer layers possess higher specific angular momentum than the inner core (see Fig. 5). As a result, angular momentum transport and centrifugal support push material outward, leading to the formation of a “centrifugal barrier”, a region where the outward centrifugal force balances gravitational attraction. As the merger proceeds, shock waves are generated as a result of the collision and compression of matter. These shocks heat the material they pass through, and as they propagate outward, they reach the regions where matter is accumulating due to the centrifugal barrier. The mate-

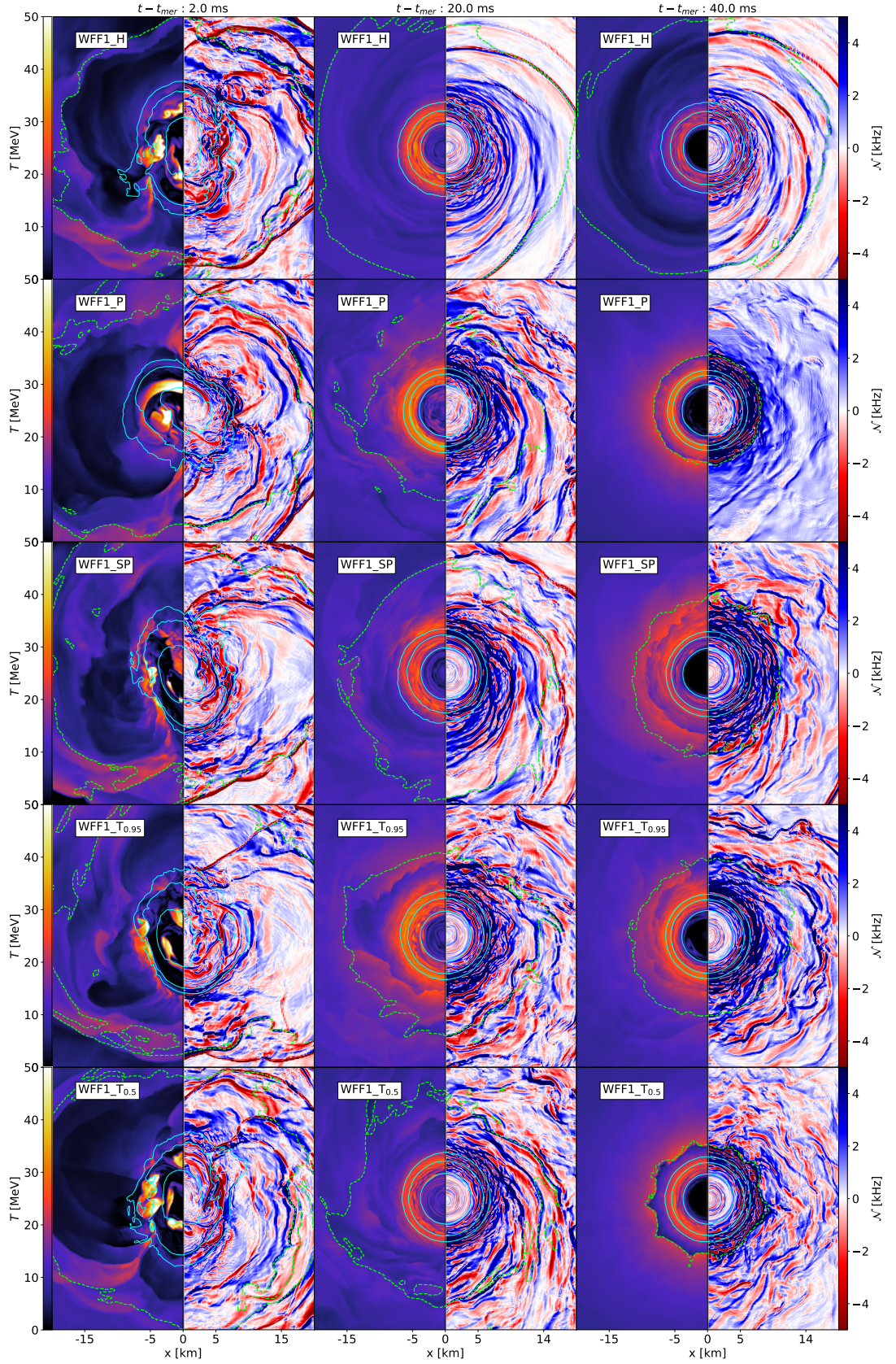


FIG. 6. Snapshots of the temperature in MeV calculated via Eq. (16) (left panels) and the Brunt-Väisälä frequency in kHz calculated via Eqs. (B1) and (B2) (right panels) for the WFF1 BNS remnants with varying magnetic field content on the equatorial plane at selected times. Cyan solid lines show different rest-mass density contours corresponding to the boundaries between segments of the piecewise equation of state used to model the star, corresponding to densities $1.76 \times 10^{14} \text{ g/cm}^3$, $5.01 \times 10^{14} \text{ g/cm}^3$, and 10^{15} g/cm^3 . The green dashed line marks the bulk of the star, defined as $10^{-3} \rho_0^{\text{max}}$.

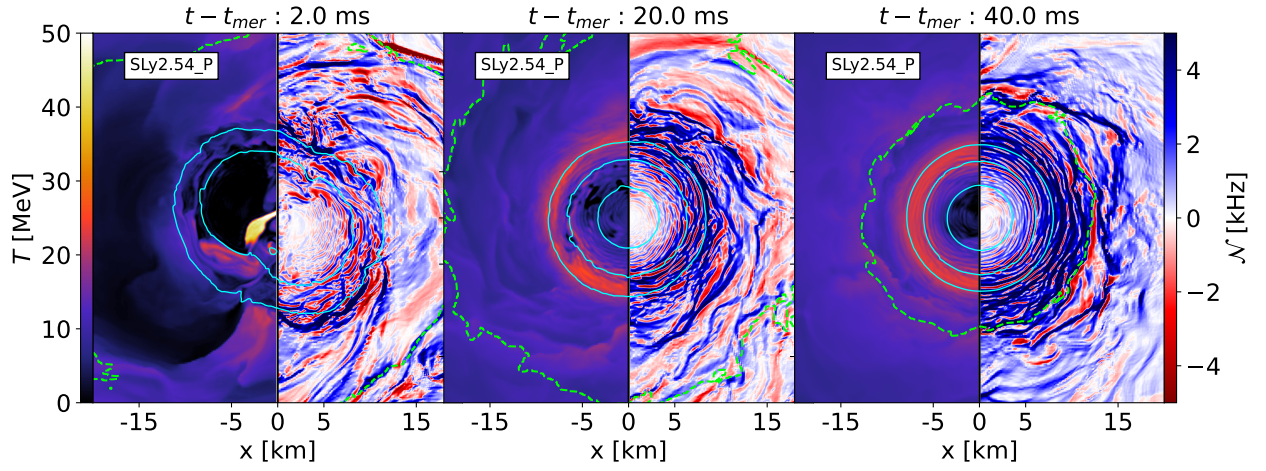


FIG. 7. Same as in Fig. 6, but for the remnant of the SLy2.54.P case. Here, cyan lines correspond to densities $1.46 \times 10^{14} \text{ g/cm}^3$, $5.01 \times 10^{14} \text{ g/cm}^3$, 10^{15} g/cm^3 .

rial in these regions reaches temperatures of $\sim 30 - 40 \text{ MeV}$. Note that in the calculation of the temperature, we do not consider the presence of thermal neutrinos. If we were to consider their presence, the temperature would be lower. Furthermore, our simulations do not include a treatment for neutrino transport, and previous simulations with such treatments have shown that the temperature decreases when introducing neutrinos [174]. Therefore, we would expect the true temperature of BNS remnants to be a few percent lower than what we report here.

In magnetized cases, magnetic instabilities amplify the magnetic field by extracting energy from differential rotation. Part of this magnetic energy is subsequently converted into heat through dissipation processes, sustaining high temperatures for longer. In contrast, the non-magnetized remnant lacks these magnetic effects and, as a result, it cools down and is the coolest configuration at $\sim 40 \text{ ms}$ after the merger (see the right column in Fig. 6).

At early times post-merger, the Brunt-Väisälä frequency shows a complex pattern of stable and unstable regions, driven by violent oscillations and inhomogeneous heating. At 20 ms after merger, the structure becomes more coherent, with extended stable regions. At 40 ms after merger, most models show reduced instability, although some cases retain persistent unstable zones.

The regions of highest temperature roughly coincide with zones of convective stability, where, although there are some small unstable (red) regions, the Brunt-Väisälä is mostly positive (see Fig. 6). In these regions, a fluid parcel that is vertically displaced will oscillate around its initial position rather than move farther away. The buoyancy forces suppress convective motions, limiting the efficiency of heat redistribution by matter. As a result, thermal energy is not efficiently transported away and accumulates, leading to a sustained higher temperature. In contrast, in convectively unstable regions, heat is more easily redistributed by fluid motions, preventing such localized heating. In Fig. 7, we show the same correlation between high-temperature regions and convectively sta-

ble regions for a different EoS, SLy, with a lower magnetic-to-gas pressure.

It has also been suggested that there is a reverse causal relationship, in which a hotter region naturally becomes convectively stable [154]. In such a scenario, increased thermal pressure in the hot region could balance buoyancy forces, suppressing convection, and supporting a stable configuration. However, in the non-magnetized case, the high-temperature ring at $t - t_{\text{merg}} = 20 \text{ ms}$ does not translate to a coincident region of convective stability 20 ms later (see first row of Fig. 6). This region does not appear to influence the long-term convective stability of the remnant, suggesting that temperature alone does not establish convective stability in this context.

On the other hand, we see that in the densest part of the star, where the temperature is the lowest, the values of the Brunt-Väisälä frequency are very close to zero, indicating that the matter is marginally buoyant. This is expected due to the small thermal and pressure gradients.

We remark the differences between the different magnetic field topologies. In particular, the pulsar-like case shows an extended and clearer stable region, with fewer red spots. Although its size and temperature profile are similar to that of configuration $\mathbf{T}_{0.5}$, the latter retains significant unstable regions, even at 40 ms after merger, consistent with sustained differential rotation, while the former pulsar-like case is almost rigidly rotating in the core (see Fig. 5). If a fluid element moves into a region with a different angular velocity, centrifugal forces can enhance the displacement rather than restore it. Note that, for a differentially rotating star, the criteria stability should account for differential rotation by also evaluating the epicyclic frequency, $\kappa^2 = 1/\varpi^3 \nabla(\Omega^2 \varpi^4) \cdot \nabla \varpi$, in which case we recover the Solberg-Høiland criteria [166]. At 40 ms after merger, cases $\mathbf{SP}$ and $\mathbf{T}_{0.5}$ retain a significant amount of differential rotation, while $\mathbf{T}_{0.95}$, although still differentially rotating, shows a less pronounced profile (see Fig. 5). If we were to account for the differential rotation, then we would expect more stability in the core of the star and less stability towards the surface.

In addition, the Brunt-Väisälä frequency provides information about the natural g-mode oscillation frequencies of the fluid in the remnant. In particular, it indicates that these frequencies are on the order of kilohertz in our systems. While for non-rotating or slowly rotating stars g-mode frequencies are typically < 0.6 kHz [175], for rapidly rotating stars they can reach the kHz range [176]. The discrepancy between the frequency range obtained in our analysis and that reported in the literature is likely due to the approximations adopted in our approach (see Appendix B).

Finally, in Fig. 6 one can also notice how the magnetic field presence and topology impacts the structure and size of the remnants. As displayed in the third column, the non-magnetized case results in a more extended remnant. This is likely due to less efficient angular momentum redistribution in the absence of magnetic fields. In magnetized cases, effective magnetic viscosity triggers outward angular momentum transport and rapid dissipation of rotational energy, allowing the remnant to contract and become more compact. In contrast, in the non-magnetized case, angular momentum is redistributed only through numerical viscosity, which is a less effective mechanism for angular momentum transport. As a result, the remnant remains more extended and supported by differential rotation (see Fig. 5). By contrast, the smallest remnants correspond to those with an initially extended poloidal component, cases **P** and **T**_{0.5}.

B. Gravitational waves

Figs. 8 and 9 show the GW strain of the dominant mode h_{22} as a function of retarded time post-merger for the cases in Table IV. Before merger, the corresponding waveforms are nearly identical across all configurations, including both magnetized and non-magnetized evolutions. This is expected, as the magnetic field is embedded in the stars two orbits prior to the merger. We note that the merger time is the same in all the cases with the same EoS and initial magnetic-to-gas pressure.

Following merger, we observe a higher-frequency signal arising from oscillations of the newly bar-deformed binary remnant. In the HMNS SLy cases, depending on the magnetic field content, and thus the effective magnetic viscosity, the post-merger GW signal decreases in amplitude and either: i) comes to an abrupt end following the quasinormal ringdown modes of the BH, as seen in the **P** and **T**_{0.95} configurations (cases with stronger magnetization; see Fig. 8), within a timescale of $t - t_{\text{merg}} \sim 3$ ms (see Table IV for details); or ii) follows an oscillatory behavior with a slowly varying envelope, as in the **H**, **SP** and **T**_{0.5} configurations, which terminates when the remnant collapses to a BH at $t - t_{\text{merg}} \sim 15$ ms.

On the other hand, the waveforms associated with the WFF1 binaries show similar behavior during the early post-merger phase, with significant differences only emerging after $t - t_{\text{merg}} \gtrsim 5$ ms (see Fig. 9). In the non-magnetized (WFF1_H) and superposition (WFF1_SP) cases, the quadrupolar oscillation modes of the fluid persist throughout the duration of the simulation, resulting in a sustained GW

signal with little post-merger damping, which is more evident in the WFF1_H than in WFF1_SP. In contrast, configurations with strong magnetic fields (i.e. WFF1_P, WFF1_T_{0.95}, and WFF1_T_{0.5}) show rapid damping of the waveform amplitude, becoming negligible by $t - t_{\text{merg}} \gtrsim 15$ ms. This suppression of post-merger oscillations indicates the formation of a uniformly rotating core, driven mainly by the MRI, which is more efficiently resolved in models with dominant poloidal fields. This interpretation is consistent with the rotational profiles (see Fig. 5), which show a magnetically braked core surrounded by a low-density, Keplerian cloud.

a. GW Spectrum and Prony's analysis for longer-lived remnants: To further analyze the behavior of the stable remnants, we study the time-frequency representation of the main GW mode. The general behavior of our four BNS systems is effectively illustrated by the spectrogram in Fig. 10, which also includes thick (purple) lines that trace the time evolution of the dominant spectral modes of the remnant. These frequencies are estimated using the ESPRIT Prony's method with a 4 ms moving window (see Sec. II G). The corresponding mode amplitudes, presented in arbitrary units, are shown in the bottom part of each panel. The relative intensity of the spectral density is also displayed, with darker areas corresponding to higher intensity. We note the spectrograms show a higher intensity shortly after merger, and not exactly at $t = t_{\text{merg}}$, as the analysis employs a moving window that captures preceding dynamics.

As the magnetic field during the inspiral epoch is only advected with the fluid, and so does not change the system, we observe that at merger the instantaneous frequency f_{max} across the WFF1 cases is 2 kHz. However, during merger magnetic instabilities are triggered, leading to a significant amplification of the magnetic field. During this epoch, all models exhibit very similar spectrograms in both frequency and time. The dominant transient f_{2i} frequency is ~ 4.0 kHz for WFF1, which coincides with twice the rotation frequency of the bar deformation of the binary remnant (see Fig. 5). However, we note a change in frequency of the order of ± 200 Hz depending of the magnetic field configuration (see Table IV), which remains throughout the post-merger window, i.e. this shift in frequency is also evident in the f_2 frequency. As shown in the bottom of each panel in Fig. 10, the amplitude of the dominant mode decreases over time, and is fully damped in SLy2.54.P and in both WFF1_P and WFF1_T_{0.95} by $t - t_{\text{merg}} \sim 25$ ms. This is somehow expected since in these cases we completely resolve the MRI (see Fig. 2). Similar shift in frequency due to magnetic viscosity have been reported in [89, 92].

Another feature similar in all the spectrograms, independently of their magnetic field topology, is the excitation of an intermittently low-frequency mode of around ~ 1.25 kHz, which may be linked to either the reduction in differential rotation (see Fig. 5), or the gradual decay of the perturbation amplitude as shown in the bottom of each panel in Fig. 10. Note that the time at which this mode is excited depends on the magnetic field configuration, and also decreases over time. Additionally, there is a third mode closer to the f_2 mode that is also active in all cases, although it stands out the most from

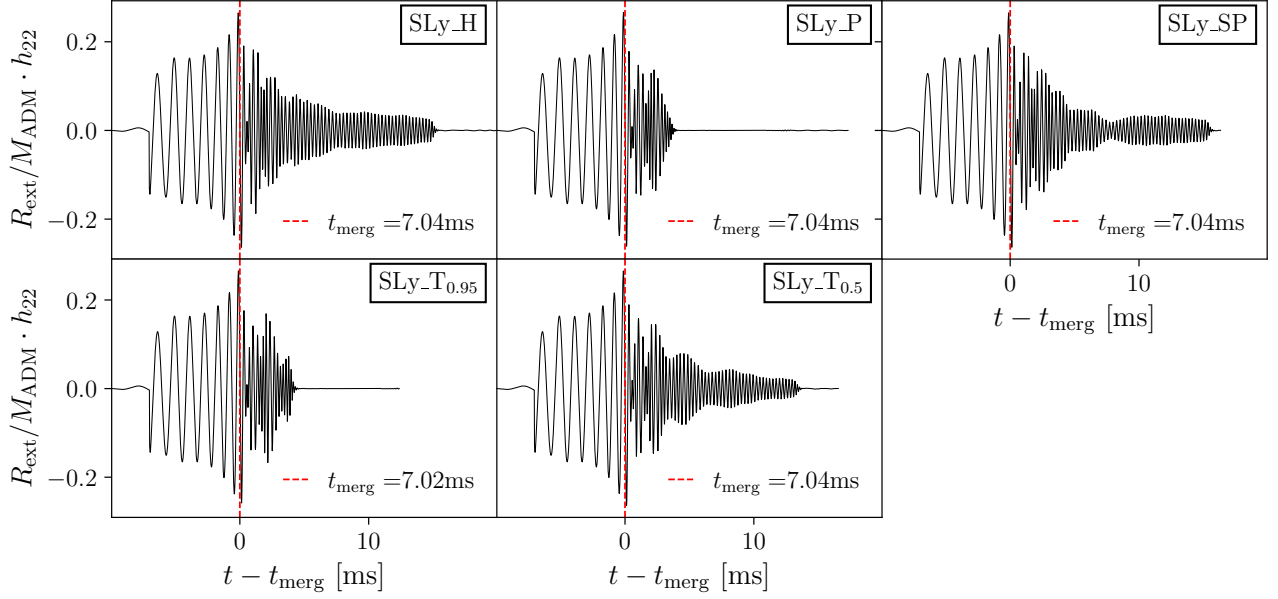


FIG. 8. Gravitational wave strain of the dominant mode h_{22} extracted at the coordinate radius $R_{\text{ext}} \sim 1240 \text{ km} \approx 840M$ for the SLy EoS cases that undergo delayed collapse to a BH (see Table IV) as a function of coordinate time measured from the merger. The merger time is defined as the peak of the GW signal that is indicated by the red dashed line.

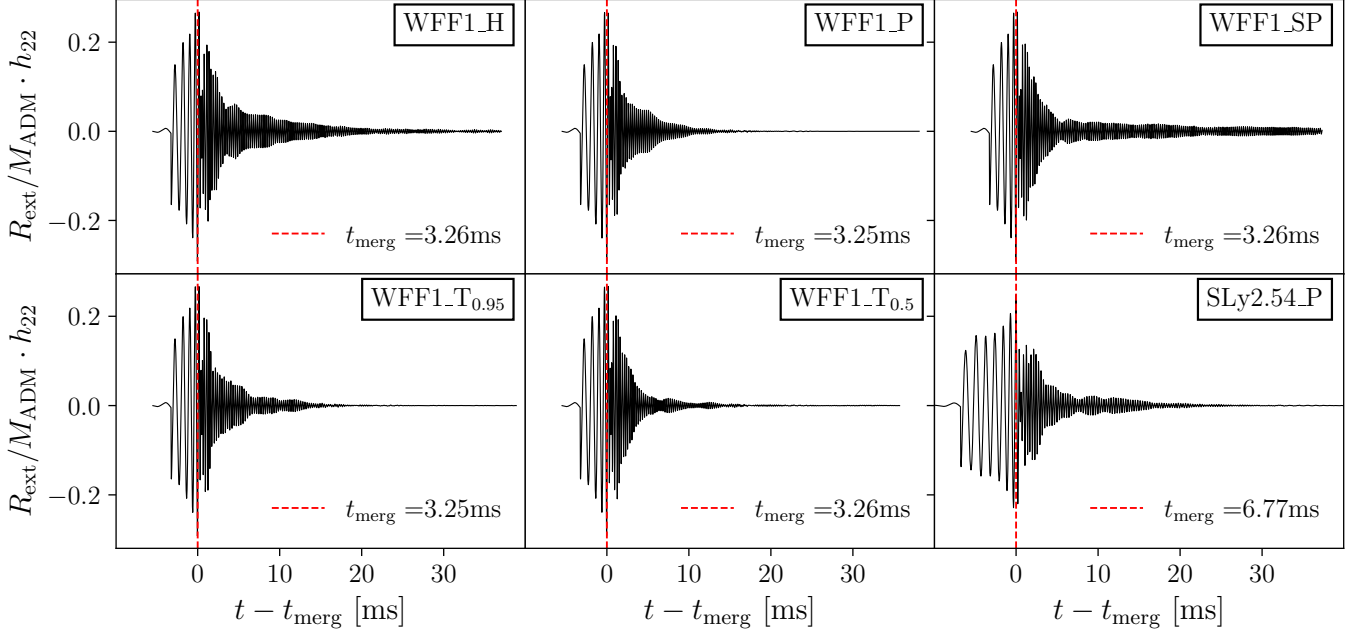


FIG. 9. Gravitational wave strain of the dominant mode h_{22} extracted at the coordinate radius $R_{\text{ext}} \sim 620 \text{ km} \approx 420M$ for binary cases that for a long-lived remnant (WFF1 EoS and SLy2.54_P in Table IV) as a function of coordinate time measured from the merger (red dashed line).

f_2 in the **P** case. The fact that the frequency of this mode is close to the frequency of the $m = 2$ mode suggests that this is not one of its overtones (see below).

b. Power spectral density (PSD): Figure 11 shows the power spectral density (PSD) of the GW signal for WFF1 cases computed using two distinct post-merger time windows.

As mentioned before, the magnetic field influences the peak frequency of the dominant mode in both the early and late post-merger phases. It is worth mentioning that the **P** model displays the highest peak frequency with $f_2 \simeq 4.11 \text{ kHz}$.

On the other hand, WFF1_SP and WFF1-T0.95 models have lower peak frequencies than that of the non-magnetized case.

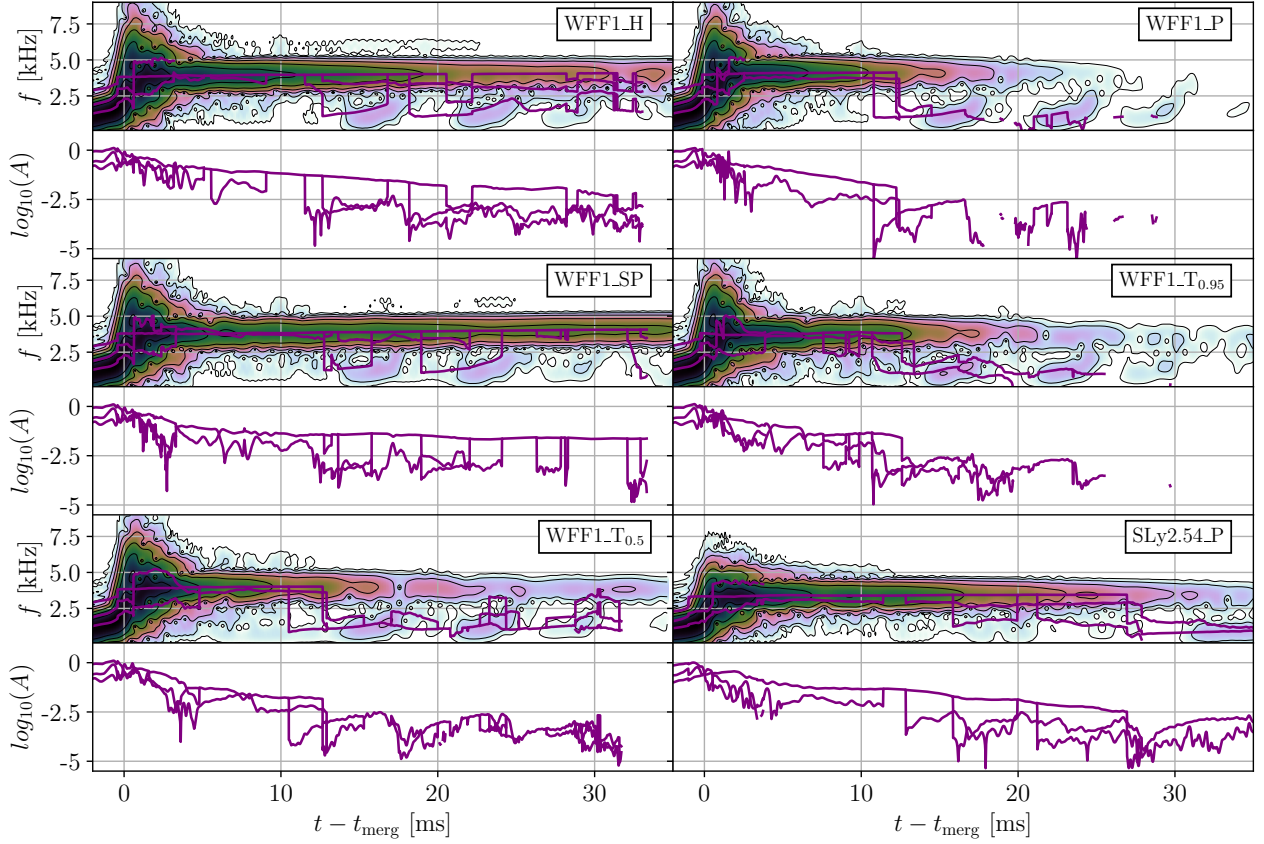


FIG. 10. The top half of each panel shows time-frequency spectrograms of the $l = m = 2$ component of the GW strain for all models. Thick purple lines, computed from the Prony analysis, indicate the frequency of the active modes responsible for most of the GW emission at different times. Colors indicate the relative intensity of the spectral density (darker areas correspond to higher intensity). The bottom half of each panel displays the amplitude of the main active modes (in arbitrary units).

As the f_2 can be correlated with the compactness of the remnant, more compact remnants tend to have higher f_2 values, while less compact ones yield lower frequencies [157, 181]. While this trend is observed for the magnetized cases, it is noteworthy that the non-magnetized remnant is actually the largest in size (see, e.g. Fig. 6, where the bulk of the star is outlined by a green dashed line). Thus, one would naively expect this case to produce the lowest f_2 frequency. The fact that configurations **SP** and **T_{0.95}** exhibit even lower f_2 values suggests that the peak frequency is influenced by other physical properties of the star remnant.

Along with the power spectral density, we also plot vertical lines corresponding to the expected values for f_{2i} and f_2 peaks from the quasi-universal relations from [181, 182]. We observe that both of these quasi-universal relations deviate for this equation of state, and that the 1σ interval fails to account for the frequency shifts caused by differing magnetic field topologies configurations. Notice that, as pointed out in [92], non-magnetized BNSs suggested that shifts in the main post-merger GW peak could be the smoking gun to feature some properties of the EoS like phase transitions. However, our results suggest that magnetic fields alone can produce similar frequency shifts, as reported in [89, 92].

c. Density eigenfunctions: To study the main oscillation modes present in the remnants, whose corresponding frequencies are shown in Fig. 10 and identified as the peaks in Fig. 11, we performed a FFT over ~ 2 ms intervals of the density distribution and identified the Fourier amplitude at selected fixed frequencies, corresponding to the eigenfunction of a particular mode as done in [154, 183, 184]. We labeled each segment using its midpoint time.

Fig. 12 shows the density eigenfunctions at the frequencies identified in the spectrograms as the main active modes, at a given time post-merger. The time displayed is different for each remnant, and is chosen based on the clearest images obtained from this analysis. Each panel is scaled to the peak value of the density eigenfunction presented.

The first and third columns display the eigenfunction of the f_2 mode for each WFF1 configuration in the equatorial plane. A telltale sign that this corresponds to the f_2 mode is its quadrupolar shape. The absence of nodal lines in this plane suggests that this mode is the fundamental one. As mentioned before, using Prony's analysis we identify a new mode with a frequency close to the fundamental one (see Fig. 10). The density eigenfunctions for this frequency are shown in the second and fourth columns. We notice WFF1.P displays a clear

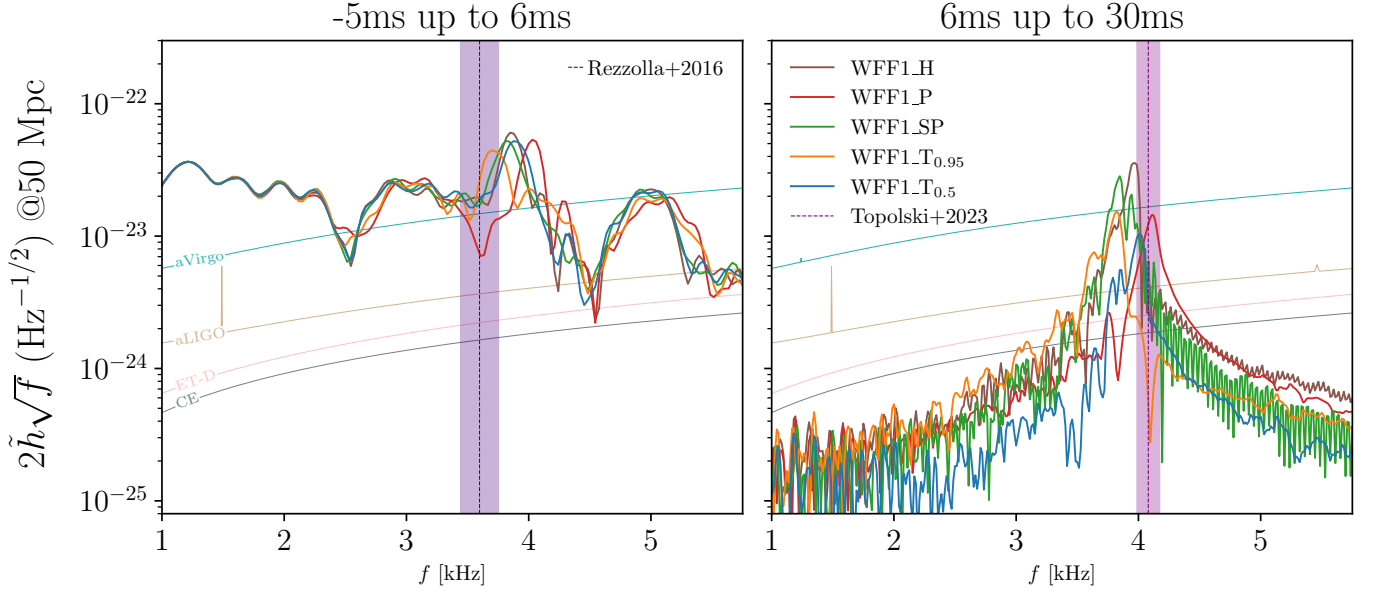


FIG. 11. GW spectra for BNS mergers with WFF1 EoS at 50 Mpc with optimal orientation. The spectra are shown for different restricted time windows (indicated at the top of each frame) in order to emphasize the contribution of the dominant spectral components at different times. For reference, we also show the sensitivity curves of Advanced Virgo [177], Advanced LIGO [178], Einstein Telescope [179], and Cosmic Explorer [180]. Vertical lines show the expected values for f_{2i} and f_2 peaks on the left and right panels, respectively, from the quasi-universal relations from [181] and [182]. The shaded region around these lines covers 1σ .

node along the radial direction. We exclude the possibility of this being an overtone of the $m = 2$ mode, since its frequency is very close to the corresponding $f_2 \approx 4.11$ kHz. However, the radial node is also a characteristic of the $m = 0$ mode (see e.g. the top panel of Fig. 1 in [184]). The frequency of the $m = 0$ mode can be extracted from the oscillations of the lapse, and it is ~ 1 kHz. Although the frequency of this additional mode is not exactly $f_2 - f_0$ or $f_2 + f_0$, we can conclude this is a non-linear coupling of modes $m = 0$ and $m = 2$ due to the topology of its density eigenfunctions, which exhibit both the quadrupolar shape of the $m = 2$ mode and the single radial node of the $m = 0$ mode. Although the additional mode of the other remnants display a similar structure, their nature is not as clear as in WFF1_P. One particular thing to look out for is the decreasing amplitude in the radial direction, which can indicate that the presence of a radial node is due to differential rotation and does not correspond to a new mode [154, 183]. Another indicator of this is that the node is farther away from the center of the star. An obvious example of this is WFF1_SP, which retains a high degree of differential rotation even at 40 ms after merger (see Fig. 5). Therefore, we can only conclude with confidence that, during the time of our evolutions, the pulsar-like configuration excites a non-linear coupling of the $m = 0$ and $m = 2$ modes.

d. Detectability of the shifts in the peak frequency:

To assess the detectability of the shifts in the peak frequency arising from different magnetic field topologies, we use the BayesWave algorithm [158], which reconstructs unmodeled GW signals using a basis of sine-Gaussian wavelets and minimal assumptions [185]. The algorithm employs

a trans-dimensional reversible jump Markov Chain Monte Carlo (RJMCMC) to jointly sample the posterior distributions of the wavelet parameters. From these, posterior distributions for derived quantities, such as the GW spectra discussed above, are obtained by applying a FFT to the reconstructed signals, with a boxcar (rectangular) window. This allows us to compute posteriors for the peak frequencies associated with the dominant post-merger modes (see discussion in Sec. III A). The analysis is performed using time-domain waveforms from WFF1 models, those that do not undergo collapse, injected into Gaussian noise consistent with the sensitivity of the Einstein Telescope (ET). We adopt the ET-D configuration [186], consisting of a three-detector network at a single site, and assume optimal sky location and inclination. For the wavelet functions, we set a maximum number of wavelets of $N_{\max} = 200$ and a maximum quality factor of $Q_{\max} = 200$. We also employ $n = 4 \times 10^6$ iterations, and a sampling rate of 16384 Hz. We use uniform offset phase parameters in the range $\phi_0 = [0, 2\pi]$, and the signal wavelet amplitude prior is presented and discussed in [187]. This setup enables us to determine the potential detectability of frequency shifts induced by varying magnetic topologies, particularly the f_2 mode and additional subdominant modes identified in the spectrograms.

The upper panel in Fig. 13 shows the mean value of the recovered peak frequencies, together with their standard deviation, as a function of the distance to the GW source. The horizontal lines represent the injected values. The time window used to compute the GW spectra is chosen to capture the early post-merger phase, i.e. from $t - t_{\text{merg}} = -1$ ms

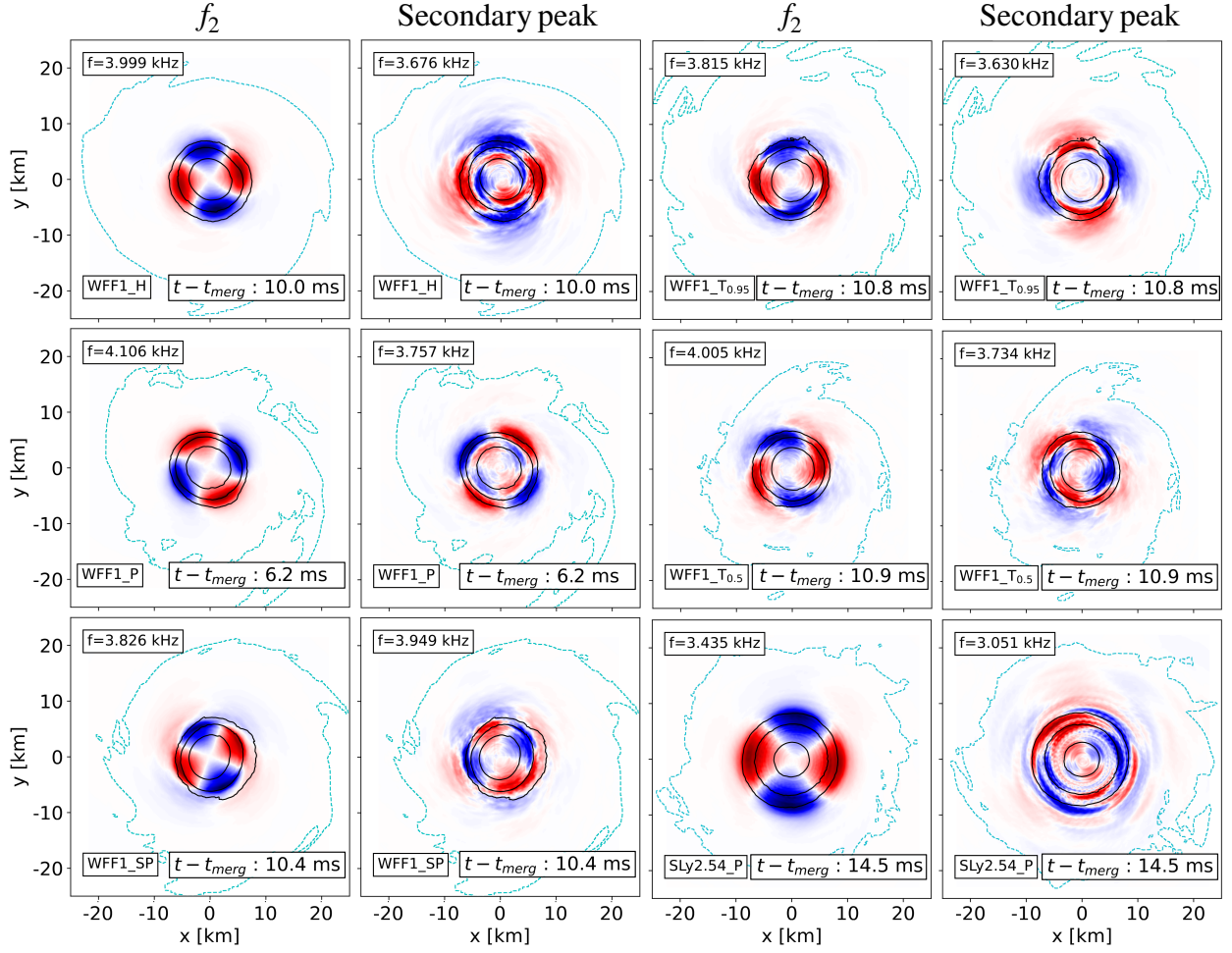


FIG. 12. Density eigenfunctions in the equatorial plane of the WFF1 models in Table IV. The black lines are isocontours of the rest-mass densities $1.76 \times 10^{14} \text{ g/cm}^3$, $5.01 \times 10^{14} \text{ g/cm}^3$, and 10^{15} g/cm^3 corresponding to the boundaries between segments of the piecewise equation of state used to model the star. The cyan line indicates the surface of the star at $10^{-3} \rho_0^{\text{max}}(t=0)$. The first and third columns display the f_2 mode, while the second and fourth show the additional mode observed in Fig. 10. The radial node in the density eigenfunctions of this mode suggests a coupling with the $m=0$ mode.

to $t - t_{\text{merg}} = 6 \text{ ms}$, where the signal amplitude is larger. For distances up to $\sim 25 \text{ Mpc}$, all the frequency peaks are distinguishable. At a distance of $D = 50 \text{ Mpc}$, most peaks are still distinguishable, but the one from WFF1_P shows a larger discrepancy, due to the peak being located at higher frequencies, where the detector noise is greater. At a distance of 75 Mpc , the recovered values overlap, and thus these shifts may no longer be detectable at this distance.

A way to check the reconstruction performance of the algorithm is by computing the overlap function between the injected and recovered signals. This function can take values from -1 to $+1$. A perfect match between the signals will result in $\mathcal{O} = +1$, and a perfect anti-correlation will give $\mathcal{O} = -1$. If there is no similarity between the signals, the overlap will be 0 . The expression for the weighted overlap of a network of

N detectors is [187]

$$\mathcal{O}_{\text{network}} = \frac{\sum_{k=1}^N \langle h_i^{(k)}, h_r^{(k)} \rangle}{\sqrt{\sum_{k=1}^N \langle h_i^{(k)}, h_i^{(k)} \rangle} \sqrt{\sum_{k=1}^N \langle h_r^{(k)}, h_r^{(k)} \rangle}}, \quad (18)$$

where h_i and h_r denote the injected and recovered signals, respectively, and the index k refers to the k -th detector. This expression involves the inner product of two complex quantities, which can be defined as

$$\langle a, b \rangle \equiv 4\text{Re} \left[\int_{f_{\min}}^{f_{\max}} \frac{a(f)b^*(f)}{S_h(f)} df \right], \quad (19)$$

where $S_h(f)$ refers to the noise PSD of the detector and $(f_{\min}, f_{\max}) = (2000, 6000)$ is the bandwidth of the analysis. Note that we only consider the post-merger frequencies of interest. Here, the asterisk indicates complex conjugation.

The lower panel of Fig. 13 shows that up to $\sim 25 \text{ Mpc}$, the network overlap is roughly 1 , implying that the injected and

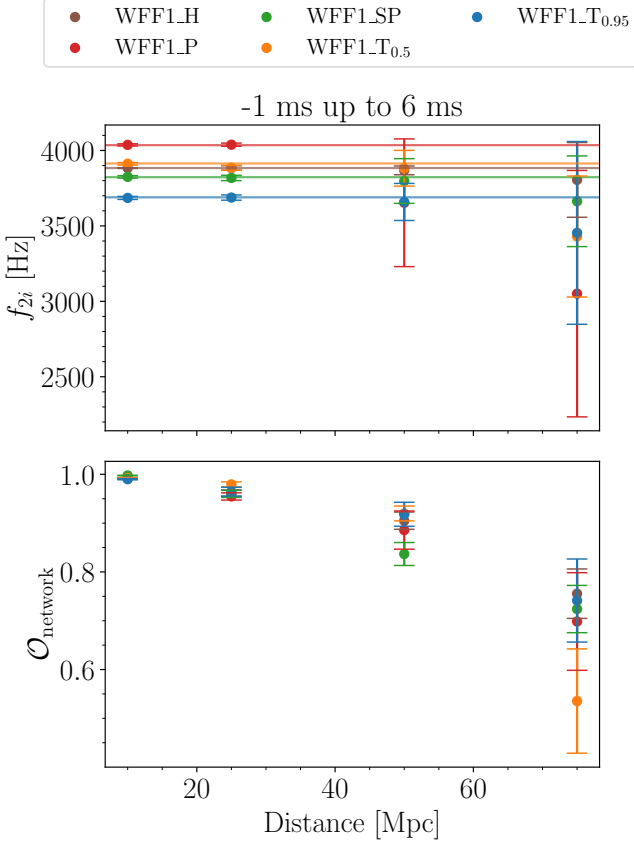


FIG. 13. Mean values (dots) of the recovered posterior distributions using the Einstein Telescope of the peak frequencies (upper panel) by the BAYESWAVE algorithm and the network overlap (lower panel), together with their standard deviation (bars), as a function to the distance to the GW source. The injected peak values are displayed by solid lines in the upper panel. The colors represent the different magnetic field topologies for the WFF1 configurations that do not undergo collapse to a BH.

recovered signals match almost perfectly. However, as the distance to the source increases, the overlap drops to lower values. At $D = 50$ Mpc, the overlap lies between $\sim 0.8 - 0.9$. The match is still large, but since the frequency peaks are very close for the different magnetic field configurations, the recovered peaks at this distance already overlap in some cases. For a distance of $D = 75$ Mpc, the network overlap falls below 0.8, in agreement with the large uncertainty in the recovered peak frequencies. Therefore, we conclude that shifts in the early f_{2i} frequency peaks arising from different magnetic field configurations could be detected by third-generation detectors, e.g. Einstein Telescope, for distances up to ≈ 50 Mpc.

C. Electromagnetic signals

The fluid and Poynting luminosities, as defined by Eqs. (11) and (12), respectively, are shown in Fig. 14 for all the cases. We notice that for both EoS the Poynting luminosity of the **SP** configuration is significantly lower than that of the other

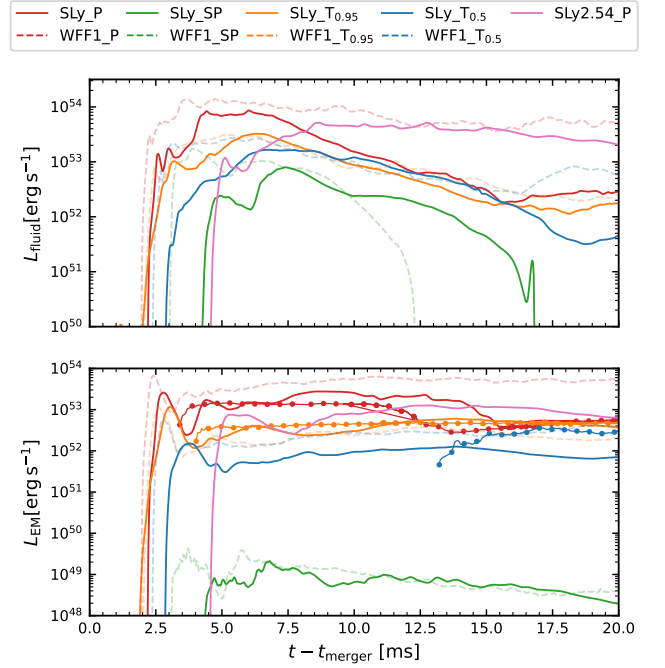


FIG. 14. Luminosity via fluid outflow (top) as computed by Eq. (11) and EM radiation (bottom) computed by Eq. (12) as a function of the coordinate time since t_{merg} (defined as the time of peak GW emission) extracted over a sphere at $r_{\text{ext}} = 539$ km. The dots and corresponding connecting lines in the Poynting luminosity plot display the expected BZ electromagnetic luminosity from Eq. (13) for the BH mass, spin and approximate B_{pol} for cases leading to a BH.

cases, with values $\sim 10^{48} - 10^{49} \text{ erg s}^{-1}$, partially consistent with the lower end of the range expected for short gamma-ray bursts ($\sim 10^{49} - 10^{53} \text{ erg s}^{-1}$) [133, 147, 148], while the other configurations are well within that range. We will see shortly that there is no evacuated funnel for the **SP** configuration. The lack of a magnetically-dominated funnel explains the fact that the fluid luminosity becomes negative for the **SP** cases around 17 and 12.5 ms after merger for the SLy and WFF1 EoS, respectively, since there is no outflow.

We note that both the fluid and Poynting luminosity of the configuration **T**_{0.95} are largely independent of the EoS (compare the solid and dashed orange lines), while case **P** shows the largest difference between the two EoS (compare the solid and dashed red lines). This difference in EoS dependence can be explained by the different outcomes of the simulations, both in terms of the fate of the remnant and the magnetization levels. The difference observed between cases SLy_P and WFF1_P is likely due to the accretion of material onto the BH in the former. This hypothesis is supported by the fact that SLy2.54.P, which does not collapse, presents fluid luminosity values of the same order of magnitude as WFF1_P. One could expect the same to be true for the SLy_T_{0.95} and WFF1_T_{0.95}, as the fate of their remnants is the same as their **P** counterparts. However, their magnetization levels are different. The second row of Fig. 15 shows the magnetization levels of the WFF1 cases at ~ 40 ms after the merger, where we see that

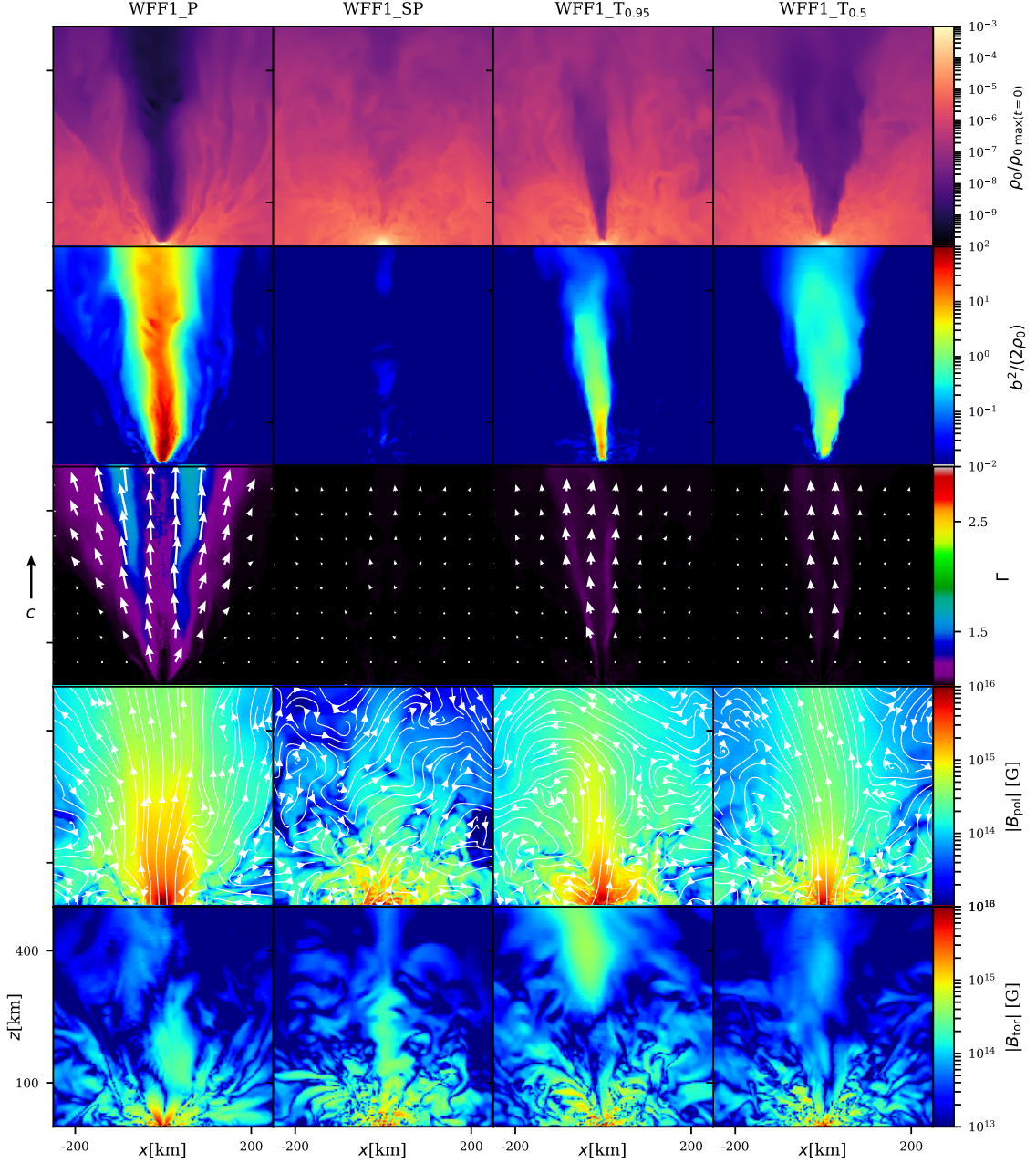


FIG. 15. Meridional slides at $t - t_{\text{merg}} \sim 40$ ms for the magnetized WFF1 BNS remnants in Table IV. We show the rest-mass density ρ_0 on a logarithmic scale relative to the initial maximum rest-mass density (top row), $b^2/2\rho_0$ which approximately corresponds to the magnetization σ (second row), the Lorentz factor (third row), and the strength of the poloidal and toroidal magnetic field (fourth and fifth rows). White lines indicate the magnetic field lines while the arrows indicate the flow velocities. The arrow labeled on the left-hand side indicates the magnitude of the speed of light.

the **P** configuration presents much larger magnetization values than the ones of $T_{0.95}$, which translates to an increased particle acceleration. The poor acceleration in the $T_{0.95}$ configuration leads to less amounts of unbound material (see Fig. 16) even if no BH was formed, and therefore, for this magnetic field configuration, the accretion onto the BH does not have a big impact on the fluid and Poynting luminosities, nor on the ejecta mass.

Unlike configuration $T_{0.95}$, and despite showing similar ejecta masses (see Fig. 16), cases with $T_{0.5}$ yield, throughout the evolution, Poynting luminosities of different orders of magnitude for each EoS. Notice that SLy_0.5 and WFF1_0.5 exhibit similar fluid luminosities, especially in the range $t - t_{\text{merg}} \sim 6.5 - 15$ ms, diverging only after SLy_0.5 collapses at $t - t_{\text{merg}} \approx 13.17$ ms (the time difference can be explained by the fact that the fluid luminosity is extracted at

$r_{\text{ext}} = 539$ km, and so the effects of the collapse can only be noticed in the fluid luminosity at later times). In particular, when SLy- $T_{0.5}$ forms a BH, the magnetic energy drops (see Fig. 4), suggesting that the acceleration of particles will be reduced for this case. However, the difference in the Poynting luminosity is roughly constant and is probably attributed to both the differences in the ejected mass as a result of the different EoS and the slight differences in the magnetic energy of these two cases.

The lower panel of Fig. 14 also displays dots with connecting lines that correspond to the expected BZ electromagnetic luminosity as computed via Eq. (13). The expected BZ luminosity for the **SP** case is below the lower limit of the y-axis range and therefore is not seen in the plot. The other three cases show good agreement between each other by the end of our simulations, in particular cases **P** and $T_{0.95}$ which collapse around the same time and whose BHs have the same mass and spin (see Table IV). The Poynting luminosity of the $T_{0.95}$ configuration agrees for the most part with the one expected from the BZ mechanism, while the pulsar-like configuration differs from the expected one around $t - t_{\text{merg}} \sim 7.5 - 15$ ms, but their agreement by the end of the simulation indicates that the BZ mechanism dominates then. In contrast, the Poynting luminosity of $T_{0.5}$ does not match the BZ one. This is likely related with the “late”-time formation of a BH in that case, which might not have allowed enough time for the BZ mechanism to become dominant by the time we terminated this simulation.

The Poynting flux generated by the rotating coiled magnetic field accelerates an outflow wind of gas, forming an evacuated funnel (see top row of Fig. 15). Naturally, this does not happen for the **SP** configuration, as the Poynting luminosity is significantly lower (there is no jet). We also note that, for the **P** configuration, the funnel has a significantly lower density than the other cases. Between configurations $T_{0.95}$ and $T_{0.5}$, the latter presents lower densities in a broader region, suggesting that a poloidal component might play a key role in evacuating the funnel. In the second row of Fig. 15, we show $b^2/(2\rho_0)$, which approximately corresponds to the magnetization, and is closely related to the terminal Lorentz factor Γ_∞ (see Eq. (8)), therefore being a good proxy for sGRB-consistent incipient jets. This is naturally connected to the first row, as lower density regions will present larger magnetization values. Thus, the pulsar-like configuration is the only one with Γ_∞ consistent with sGRBs ($\sim 10^2$). While $T_{0.95}$ reaches large Γ_∞ values, this is restricted to the funnel region close to the remnant. With smaller values of $b^2/(2\rho_0)$, $T_{0.5}$ presents larger values than $T_{0.95}$ farther away, consistent with its clearer funnel. The third row of Fig. 15 shows that the **P** case accelerates the fluid more efficiently, reaching Lorentz factors within the simulation box of > 1.5 , whereas case $T_{0.95}$ reaches at most $\Gamma \sim 1.3$, and the other cases presenting even smaller values. The larger values in case $T_{0.95}$, when compared to $T_{0.5}$, can be explained by poloidal magnetic field strength in these cases (see the fourth row of Fig. 15). While $T_{0.5}$ ’s magnetic field lines seem to be more collimated, the poloidal component is stronger in $T_{0.95}$, leading to larger Lorentz factors.

We note that, although case $T_{0.95}$ starts with an exterior

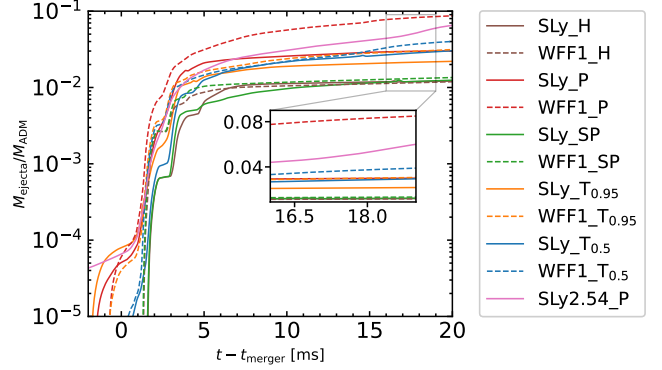


FIG. 16. Rest-mass fraction of the ejected mass M_{ejecta} following the GW peak amplitude computed in the domain $r > 30M$.

poloidal field, its “late”-time exterior topology differs significantly from that of case **P**. This is not unexpected as most of the magnetic energy is concentrated in the stars. However, we do note that there is some degree of collimation, especially until ~ 350 km above the equatorial plane (see Fig. 1 and the fourth row of Fig. 15).

On the other hand, the **SP** case shows no degree of collimation, and the strength of the poloidal component of its magnetic field is considerably lower than that of the other cases, which we already had observed upon magnetic field insertion (see Table III) and also expected as we do not completely capture the MRI.

Consistent with Fig. 4, the pulsar-like configuration presents both the strongest poloidal and toroidal magnetic field components. We conclude that the only topology (from the ones tested) compatible with a highly collimated, relativistic outflow of matter consistent with sGRBs is a purely poloidal one.

1. Ejecta mass and kilonova estimates

Figure 16 shows the dynamically ejecta rest-mass fraction of the unbound material computed via Eq. (10). We note that depending on the magnetic field topology the ejecta ranges from $\sim 0.5\%$ and $\sim 3.0\%$ of the total mass of the system, corresponding to $M_{\text{eje}} \sim 10^{-2} - 10^{-1} M_\odot$ (see Table IV), with average velocity $\langle v_{\text{eje}} \rangle \sim 0.2 - 0.3 c$. Ejecta masses $\gtrsim 10^{-3} M_\odot$ are expected to lead to detectable, transient kilonova signatures [188] originating from the decay of radioactive nuclei synthesized via the r-process within the neutron-rich material ejected during BNS mergers.

We note that, regardless of the magnetic field configurations, the dynamical ejection of matter for the SLy binaries begins ~ 1 ms before that from the WFF1 ones. This is expected, since the SLy binaries are less compact than WFF1 ones, and therefore, they experience stronger tidal deformation before merger. This deformation stretches the outer layers even before contact, priming them for earlier dynamical ejection. However, we observe that near the termination of our simulations, the soft-WFF1-EoS binaries exhibit a larger

overall ejection of material. This may be attributed to the higher compactness of WFF1 stars, which leads to a more violent merger interface, driving enhanced mass ejection at later times.

On the other hand, we observe that the dynamical ejection of matter in binaries with $\mathbf{T}_{0.95}$, begin slightly early than the $\mathbf{P}$ configurations. Toroidal fields, being confined to the stellar interior, enhance local magnetic pressure and promote stronger compression and shock heating during merger [39]. Additionally, they impose less magnetic tension opposing radial outflows compared to poloidal fields, allowing earlier development of dynamical ejecta [189].

To gain insight into the possible EM signatures from the dynamical ejection of matter, we employ the analytical model presented in [190]. This model computes the peak rise times, bolometric luminosities and the effective temperatures of the kilonova assuming that the ejecta is spherically distributed and expanding homologously with an average speed v_{eje} and characterized by a gray opacity κ_γ .

The peak time of the transient emission can be estimated as

$$\begin{aligned} \tau_{\text{peak}} &\sim \sqrt{\frac{M_{\text{eje}} \kappa_\gamma}{4\pi \langle v_{\text{eje}} \rangle c}} \\ &\approx 4.6 \text{ days} \left(\frac{M_{\text{eje}}}{10^{-2} M_\odot} \right)^{1/2} \left(\frac{\langle v_{\text{eje}} \rangle}{0.1c} \right)^{-1/2}. \end{aligned} \quad (20)$$

We assume a high opacity $\kappa_\gamma = 10 \text{ cm}^2 \text{ g}^{-1}$, which corresponds to ejecta containing an electron fraction $Y_e \lesssim 0.25$, a typically value expected from the early post-merger phase in BNSs [191]. The peak luminosity can be computed as

$$L_{\text{knova}} \sim 2.4 \times 10^{40} \left(\frac{M_{\text{eje}}}{0.01 M_\odot} \right)^{0.35} \left(\frac{\langle v_{\text{eje}} \rangle}{0.1c} \right)^{0.65} \text{ erg s}^{-1}. \quad (21)$$

Finally, assuming black body emission, and using the Stefan-Boltzmann law, the effective temperature at the peak can be estimated as

$$T_{\text{peak}} \sim 2.15 \times 10^3 \left(\frac{M_{\text{eje}}}{0.01 M_\odot} \right)^{-0.16} \left(\frac{\langle v_{\text{eje}} \rangle}{0.1c} \right)^{-0.09} \text{ K}. \quad (22)$$

While simplifying, these assumptions provide a useful first-order model that generally fits observed data. For our simulations, we find that the bolometric luminosity is $L_{\text{knova}} \sim 10^{40.6 \pm 0.1} \text{ erg s}^{-1}$ with rise times of $\tau_{\text{peak}} \sim 2.9 - 9.4$ days and an effective temperature of $T_{\text{peak}} \sim 10^{3.3} \text{ K}$, which is consistent with those obtained from other numerical simulations [51, 57] (see [188] for a review). This temperature can be converted into a peak wavelength $\lambda_{\text{peak}} \sim 1.35 \times 10^3 \text{ nm} (T_{\text{peak}}/10^{3.3} \text{ K})^{-1}$. Therefore, $\lambda_{\text{peak}} \sim 730 - 1830 \text{ nm}$, and the emission can be detected with current or planned telescopes, such as ALMA or the Vera C. Rubin observatory [192]. As expected, we find that EM signatures (Poynting and kilonova luminosities) depend more heavily on the magnetic field topology than on the EoS.

IV. CONCLUSIONS

Multimessenger detections of electromagnetic signals and gravitational waves, such as GW170817 and its EM counterparts, have proven crucial in improving our understanding of the sources of these signals. Future detections will allow us to further constrain the EoS of neutron stars and to understand their magnetic fields. With the improved sensitivity of future detectors in sight, we expect an increased quality and quantity of data to study these systems. Therefore, it is of the utmost importance to model these events beforehand, in order to accurately recover the properties of their sources.

The interior magnetic field of neutron stars is, in particular, one of the big questions surrounding these objects. Motivated by observations of pulsars' magnetic fields, GRMHD simulations of BNS mergers typically consider a dipole pulsar-like structure for the magnetic field, usually confined to the interior of the stars. However, their interior field remains unknown. In this study, we consider four distinct magnetic field topologies: a dipole pulsar-like field ($\mathbf{P}$), a linear superposition of poloidal and toroidal components inside the star ($\mathbf{SP}$), and a confined tilted toroidal component transitioning to a pulsar-like topology at radii $0.95 R_{\text{NS}}$ ($\mathbf{T}_{0.95}$) and $0.5 R_{\text{NS}}$ ($\mathbf{T}_{0.5}$). We analyze their evolution, EM and GW signals, as well as, if applicable, their remnant properties. We consider two different equations of state, SLy and WFF1, and fix the maximum value of $\beta^{-1} \equiv P_{\text{mag}}/P_{\text{gas}} = 0.003125$ inside the stars. Additionally, we consider a lower-mass system with the SLy EoS, a pulsar-like field and $\beta^{-1} = 0.0023$, in order to confirm the same behavior we observe in the WFF1 remnants, as all other SLy simulations lead to a black hole.

We perform an extensive multimessenger analysis of the emission properties of the systems, both gravitational waves and electromagnetic signatures. We study the main active modes that contribute to gravitational emission and their corresponding density eigenfunctions. We further analyze the properties of the remnants, namely their rotational profile, temperature distribution, and convective stability.

Our main results can be summarized as follows:

- Magnetic field topology significantly impacts GW emission, altering both the collapse time and the f_2 peak frequency in longer-lived remnants;
- Frequency shifts in the early f_{2i} mode due to different magnetic fields will be detectable with third-generation detectors up to $\sim 50 \text{ Mpc}$;
- Longer-lived remnants exhibit additional frequency modes beyond the fundamental $l = m = 2$, some identified through an extra radial node in their eigenfunctions. In the pulsar-like case, this mode is consistent with a non-linear coupling between the $m = 0$ and $m = 2$ modes;
- All configurations produce Poynting luminosities consistent with sGRBs, and are compatible with kilonova emission;

- A purely poloidal configuration is the most efficient in launching a relativistic jet;
- All cases are compatible with kilonova emission, regardless of magnetic field topology;
- Hot rings form at the centrifugal barrier and correlate with regions of convective stability.

Finally, we stress that future work should incorporate neutrino transport, more advanced magnetic topologies, and finite-temperature EoSs to refine the multimessenger predictions and clarify the role of NSs as potential sGRB progenitors.

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Appendix A: Magnetic Power Spectrum

To quantify the development of magnetohydrodynamic turbulence in the post-merger remnant, we compute the mag-

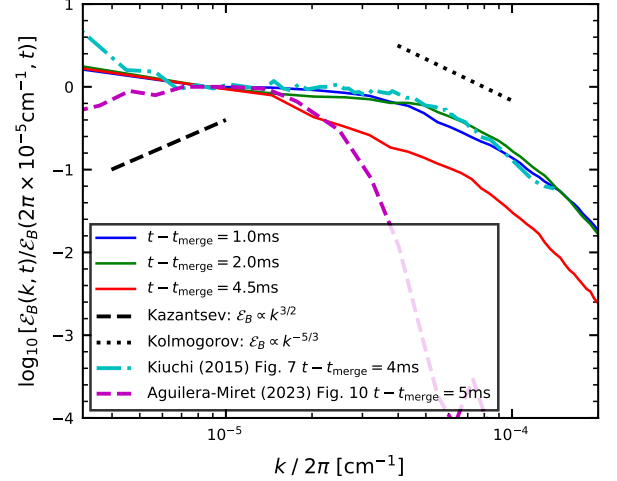


FIG. 17. Magnetic power spectrum computed via [43] Eq. B1, as used in [47], at different times post-merger for the WFF1_P simulation.

netic power spectrum following Eq. (B1) of [43]. The magnetic power spectrum $\varepsilon_B(k, t)$ characterizes the distribution of magnetic energy across spatial scales and provides insight into the dominant physical processes driving magnetic field amplification and energy cascade. Specifically, it allows us to assess whether the post-merger magnetic field growth is governed by a small-scale dynamo (e.g., Kazantsev spectrum, $\varepsilon_B \propto k^{3/2}$) or fully developed turbulence (e.g., Kolmogorov spectrum, $\varepsilon_B \propto k^{-5/3}$). We present the results for two specific simulations, involving the two different EoSs and two distinct magnetic field configurations. In Fig. 17, the magnetic power spectrum of WFF1_P is consistent with previous high-resolution studies [40], especially for earlier times post-merger. At smaller scales ($k/2\pi \gtrsim 7 \times 10^{-4}$), the spectra resemble a Kolmogorov-like slope, and at later times they do so for $k/2\pi \gtrsim 10^{-5}$, indicating a transition toward turbulent energy redistribution across scales. In Fig. 18, the magnetic power spectrum of SLy_SP shows growth, for $k/2\pi \lesssim 9 \times 10^{-4}$, with slopes compatible with the Kazantsev regime, suggesting amplification via small-scale dynamo action. The precise shape and evolution of the power spectrum remain sensitive to numerical resolution, filtering methods used to compute the Fourier transform, and the physical parameters of the system. Further studies are needed to robustly determine the spectrum’s behavior and its dependence on such parameters. Nevertheless, our results are in broad agreement with higher-resolution and large-eddy studies [40, 47].

Appendix B: Convective Stability

The Brunt-Väisälä frequency $\mathcal{N}$ quantifies the oscillation frequency of a displaced fluid element around its equilibrium position [165, 166]. It is closely associated with the Ledoux

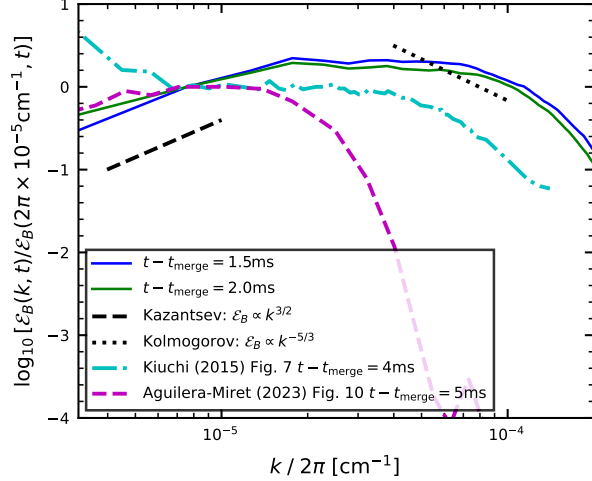


FIG. 18. Magnetic power spectrum computed via [43] Eq. B1, as used in [47], at different times post-merger for the SLy_SP simulation.

discriminant $\mathcal{B}_i$, whose sign indicates the convective stability of a given region [165, 193]. Therefore, we can compute $\mathcal{N}^2$ to understand what regions are convectively stable/unstable. In particular, a region with a positive $\mathcal{N}^2$ is stable against convection and propagation of g-modes. Furthermore, it provides an estimate of the frequency of the g-modes. We have that $\mathcal{N}^2 \propto \mathcal{B}_i \mathcal{G}_i$ [167, 169, 194], where $\mathcal{B}_i$ is the relativistic version of the Ledoux discriminant, and $\mathcal{G}_i$ is the gravitational acceleration. The proportionality factor is coordinate dependent. In the conformally flat condition approximation, this factor is α^2/ψ^4 [194], where α is the lapse function and ψ the conformal factor. For neutron stars, both α and ψ are of the order units, as can be seen from post-Newtonian analysis [195].

On the other hand, considering the metric for an axisymmetric rotating star, as was recently done in [167], this factor depends on γ_{rr} , and more generally on g_{rr} . In either case, the proportionality factor is expected to be $\sim \mathcal{O}(1)$, which differs the most from unity at the core where the gravitational field is the strongest. We have

$$\mathcal{B}_i \equiv \frac{\partial_i e}{\rho_0 h} - \frac{1}{\Gamma_1} \frac{\partial_i P}{P}, \quad (\text{B1})$$

with e the energy density, h the specific enthalpy, Γ_1 the adiabatic index around a pseudo-barotropic equilibrium, ρ_0 the rest-mass density and P the pressure. Additionally, the gravitational acceleration is given by $\mathcal{G}_i = -\partial_i \ln \alpha$, with α the lapse function. In hydrostatic equilibrium we can also write [194]

$$\mathcal{G}_i \equiv -\partial_i \ln \alpha = \frac{1}{\rho_0 h} \partial_i P. \quad (\text{B2})$$

We compute $\mathcal{N}^2$ in the equatorial plane using cylindrical coordinates. Simulations of core-collapse supernovae have shown that the two different definitions of $\mathcal{G}_i$ in Eq. (B2) differ at most by a factor of 2 [194] (see the lower right panel of their Figure 11). In Fig. 6 and 7, we compute $\mathcal{N}$ via

$$\mathcal{N} = \mathcal{N}^2/|\mathcal{N}|, \quad \text{with} \quad (\text{B3})$$

$$\mathcal{N}^2 \approx \left(\frac{\partial_r e}{\rho_0 h} - \frac{1}{\Gamma_1} \frac{\partial_r P}{P} \right) \frac{\partial_r P}{\rho_0 h}. \quad (\text{B4})$$

We can use it as a proxy for convective stability. However, its values should be interpreted with caution, as they represent only rough estimates of the true value due to the approximations employed in this work. Furthermore, a complete study of convective stability in magnetized remnants should consider dynamo mechanisms, such as the α - Ω -dynamo and Tayler-Spruit dynamo. We remark that a detailed study of convective stability is beyond the scope of this work.

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