

Magnetic field configurations in binary neutron star mergers. II. Inspiral, merger and ejecta

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(Received 31 August 2025; accepted 25 August 2026; published 24 September 2026)

We perform a series of simulations of magnetized binary neutron star mergers, with different initial magnetic field topologies, as well as varying equations of state, and mass ratios. In this paper, a companion paper to E. M. Gutiérrez, W. Cook, D. Radice, S. Bernuzzi, J. Fields, P. Hammond, B. Daszuta, H. Bandyopadhyay, and M. Jacobi, companion paper, *Phys. Rev. D* **114**, 063049 (2026), we analyze the impact of the initial field configuration on the emitted gravitational wave signal, the global amplification of the magnetic field, and the nature of the ejected material from the binary. We investigate the dependence of the phase evolution of the gravitational wave signal in the postmerger on the initial magnetic field, finding that dephasing effects between the dominant ($\ell = 2, m = 2$) mode of the gravitational wave signal and leading subdominant modes (2, 1), (3, 3) may be strongly impacted by the choice of numerical reconstruction scheme. The magnetic field amplification at merger may be considerably enhanced during the Kelvin-Helmholtz instability dominated phase by the presence of antialigned fields between the stars, or suppressed by the presence of toroidal fields, while the postmerger amplification of the field due to winding may also be suppressed by toroidal fields, and may be enhanced by asymmetries or mixtures of poloidal and toroidal fields. The magnetic field strength in the ejecta may also be strongly impacted by the nature of the initial magnetic field configuration, with initial data configurations which lead to large amplifications and those with mixtures of poloidal and toroidal fields preferentially emitting highly magnetized material in the polar regions, while we find that the magnetic field strength in the ejecta for these configurations shows a weaker dependence on the density of the material than in those cases that amplify the magnetic field less. We investigate the orientation of the field in the ejecta, and find that the magnetic field is largely randomly oriented in the ejected material, providing supporting evidence for such models used to estimate thermalization timescales of ejected material. Further we find that configurations which begin with an initial bitant symmetry break this symmetry in a uniform manner, independent of the initial magnetic field configuration, when evolved without an enforced symmetry. This behavior suggests the presence of a spontaneous symmetry breaking bifurcation in the solution.

DOI: [10.1103/1n82-rlhp](https://doi.org/10.1103/1n82-rlhp)

I. INTRODUCTION

Neutron stars (NSs) can be the site of some of the strongest magnetic fields in the Universe [1,2], reaching strengths of up to 10^{16} G. Consequently, magnetic fields

can play a key role in the evolution of binary neutron star (BNS) systems, both during the moment of merger, and in the postmerger phase. At the moment of first contact between two NSs, it is expected that a thin shearing layer will form, within which the Kelvin-Helmholtz instability (KHI) will be triggered, forming small vortex like structures in the fluid, which wind the magnetic field with them, amplifying the magnetic field energy by at least 6 orders of magnitude [3–5]. Strong magnetic fields can lead to additional pressure and tension forces in the fluid evolution which can drive the ejection of neutron rich material during

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a BNS merger [6–15], material which subsequently will undergo r-process nucleosynthesis [16,17], leading to the formation of heavy elements, visible in a kilonova [18–24].

Strong magnetic fields in the postmerger remnant have been proposed as a mechanism for the launching of ultrarelativistic jets [25–29], which are believed to lead to the production of gamma ray bursts such as GRB170817, seen coincidentally with the gravitational wave (GW) signal from the BNS merger GW170817 [30–32]. Such jets may be launched either from a highly magnetized remnant NS, a so-called magnetar, or from a magnetized black hole with an associated accretion disk, through the Blandford-Znajek process, resulting from gravitational collapse after the BNS merger [33–35].

Observations of pulsars suggest a large scale dipolar field [36–41] with contributions from higher multipoles, as the dominant exterior contribution to the magnetic field in a neutron star, developing both during the formation of the NS and through convective and thermal processes occurring in the early NS evolution [42]. The details of the interior magnetic field of the star are, however, less well constrained. Long term simulations of static neutron stars with initially purely poloidal fields [43–51] suggest that well-known magnetohydrodynamical (MHD) instabilities associated to $m = 0$ and $m = 1$ mode instabilities of poloidal configurations are important in the initial rearrangement of the field, developing a nontrivial toroidal component on Alfvén timescales [52–58], the characteristic timescale of the magnetic field evolution, $\tau_A = R_{\text{NS}} \sqrt{4\pi\rho}/B$. Magnetized BNS simulations are, however, frequently initialized with purely poloidal initial field profiles, with recent exceptions [59–63]. For binary systems which have evolved over many Alfvén timescales, we can view the constituent stars as effectively isolated from their binary partner from the perspective of their electromagnetic and thermodynamical evolution, before they approach sufficiently close through GW emission that the magnetospheres of the stars intersect, or for tidal effects to become appreciable in magnitude. During this effectively isolated evolution, we can expect that such initial linear instabilities, which operate over Alfvén timescales, have had sufficient time to saturate and that a full nonlinear rearrangement of the magnetic field has occurred. Equilibrium configurations of magnetic fields inside neutron stars have been constructed with incredibly strong toroidal magnetic fields [45], and some numerical studies have suggested that the presence of an initially strong toroidal magnetic field may be preserved over multiple Alfvén timescales [49,64]. If such configurations are indeed dynamically stable, then they may be found in the interior of stars in binary systems.

Recent works have investigated the impact of magnetic field topology on BNS evolutions. In [62,63] the impact on the postmerger gravitational wave signal has been investigated, finding that the configuration of the magnetic field may have a detectable impact on the peak postmerger

frequency. Magnetic field configurations have also been compared in [61,65,66], with results suggesting that the end state of the merger process results in a universal magnetic field configuration, independent of the initial data. Contrastingly, simulations performed in [67] have observed different magnetic field amplifications for a magnetic field localized in the crustal region of the star. Beyond magnetic field topologies, the results of [68] have suggested further universality in the extent of the amplification of magnetic fields, independent of the mass ratios, and equations of state (EOS) considered. In [59] the impact of the magnetic field orientation on jet launching is investigated, finding that the formation of a large scale field is independent of the initial condition, while [60] finds that the jet launching time, and ejecta, may be very sensitive to the specific magnetic field configuration.

The nature of the magnetic field within the ejected material from the binary may play an important role in understanding its evolution, and the associated kilonova. It has been suggested that strong magnetic fields in the ejected matter can affect the r-process nucleosynthesis, through the quantization of electron momentum states into Landau levels, affecting β - and neutron decay rates in the ejected material [69–71]. In addition the nature of the magnetic field configuration in the ejected material may affect the ability of the ejected material to fully thermalize, having a strong impact on the luminosity of kilonova light curves [72].

This paper serves as a companion paper to [73] (henceforth Paper I). We perform and discuss the results of a sequence of simulations, investigating different magnetic field configurations, different symmetry considerations, different equations of state, and different mass ratios. In particular, we perform simulations of poloidal magnetic field configurations with different geometries that are misaligned with each other to investigate generic configurations. We also perform simulations of magnetic field configurations with a combination of poloidal and toroidal fields which aim to reproduce the effect of the saturation of the initial instabilities of a purely poloidal magnetic field profile, approximating configurations found in [51]. In Paper I we focus on the resultant evolution of the postmerger remnant and disk. In this paper we focus on the gravitational wave signal, global magnetic field dynamics at merger, and the evolution of the ejected matter. We perform simulations in this paper using the numerical relativity (NR) code GR-Athena++ [50,74–76], based on the Athena++ infrastructure [77].

In Sec. II we discuss the initial data configurations we consider, and the numerical approach taken in our simulations. In Sec. III we discuss the results of our simulations, focussing on the GWs (Sec. III A), the magnetic field amplification (Sec. III B), the symmetry of the configurations found (Sec. III C) and the ejected material from the binary (Sec. III D). We summarize our findings in Sec. IV. In the

TABLE I. Summary of runs performed. B_0 represents the maximum magnetic field strength. Recon. describes the reconstruction scheme used.

Run name	B_0	Resolution	EOS	Recon.	Field configuration and comments
POL-UU	5×10^{15} G	SR	SFHo	WENOZ	Aligned poloidal
HR-POL-UU	5×10^{15} G	HR	SFHo	WENOZ	Aligned poloidal
MIX	5×10^{15} G	SR	SFHo	WENOZ	Superposed poloidal and toroidal
HR-MIX	5×10^{15} G	HR	SFHo	WENOZ	Superposed poloidal and toroidal
TOR-PP	5×10^{15} G	SR	SFHo	WENOZ	Aligned toroidal
TOR-PM	5×10^{15} G	SR	SFHo	WENOZ	Antialigned toroidal
B0	0G	SR	SFHo	WENOZ	No field
BITANT	5×10^{15} G	SR	SFHo	PPM	Aligned poloidal with bitant symmetry
LOWB	5×10^8 G	SR	SFHo	PPM	Aligned poloidal
POL-UD	5×10^{15} G	SR	SFHo	PPM	Antialigned poloidal
POL-90D	5×10^{15} G	SR	SFHo	PPM	Poloidal in orbital plane
POL-ASYM	5×10^{15} G	SR	SFHo	PPM	Displaced aligned poloidal
STIFF	5×10^{15} G	SR	DD2	PPM	Aligned poloidal
POL-Q12	5×10^{15} G	SR	SFHo	PPM	Aligned poloidal, mass ratio $q = 1.2$

Appendix, we discuss the errors associated to the numerical reconstruction procedure present in our results. Unless stated otherwise, our results are presented in centimeter-gram-second units.

II. METHOD

We perform a series of magnetized binary merger simulations using the GR-Athena++ code, with varying initial magnetic field configurations. We refer readers to Table I in Paper I, which details the overall parameters of the configurations considered, and below recap the key details of the simulations performed, which we also demonstrate in Table I.

A. Numerical details

GR-Athena++ employs task based parallelism and an oct-tree, block based grid structure, to ensure highly scaling performance on large numbers of CPU cores, as described in [77]. The Einstein equations are evolved using the Z4c formulation [78–81], with the moving puncture gauge [82–84], using a vertex centered grid and sixth order finite differencing stencils. For discussion of the vertex centered implementation we refer the reader to [74,76]. We refer the reader to [50] for the choice of gauge, constraint damping and dissipation parameters, which remain identical between all runs considered in this paper and the BNS simulations presented there.

The relativistic Euler equations are evolved in the standard $3 + 1$ “Valencia” formulation, with the Maxwell equations included in the ideal MHD approximation [85–94]. Evolution of the relativistic Euler equations is performed using a Godunov style finite volume approach, with

the Local Lax Friedrichs approximate Riemann solver used to calculate a numerical flux function, and primitives reconstructed to the cell interface with either the piecewise parabolic method (PPM) [95,96] or weighted essentially non-oscillatory Z (WENOZ) [97] stencil for high resolution shock capturing (see Appendix). In addition to the Euler equations we evolve the electron fraction Y_e as a scalar advected with the fluid density as described in [50]. The equation of state is given by tabulated data for the SFHo [98] and DD2 [99] equations of state from the Compose database [100]. External to the evolved stars we set a low density unphysical atmosphere to avoid excessively sharp gradients at the star surface, and to avoid the breakdown of the weak solutions to the Euler equations in vacuum. The density of this atmosphere is set to $\rho_{\text{atm}} = 1.8 \times 10^3$ g/cm³ and atmosphere temperature to 0.1 MeV. The electron fraction of the atmosphere is 0.5. The computational grids employed have outer boundaries set at ± 1511 km, with the base level grid sampled with 128^3 points. On top of this base level, grid refinement is added, with our standard resolution (SR) simulations adding 7 levels, giving a grid spacing of 180m on the finest refinement level fully covering the stars. We also perform two shorter-lived high resolution (HR) configurations with a finest resolution of 123m, given by a base level sampling of 192^3 points and 7 refinement levels.

The magnetic field is evolved using the constrained transport algorithm implementation of [101,102], with the divergence and curl preserving restriction and prolongation operators of [103], utilizing a face averaged sampling of the magnetic field. This provides conservation of the magnetic field divergence to machine precision, as demonstrated in [50]. Time evolution is performed using a third order

Runge-Kutta time integrator with Courant-Friedrichs-Lewy factor 0.25. Conserved-to-primitive variable inversion is implemented through a custom solver “PrimitiveSolver,” detailed in Appendix A of [50], based on the RePrimAnd algorithm of [104]. We extract gravitational waves using the Newman-Penrose scalar Ψ_4 as sampled on geodesic spheres to avoid coordinate singularities and sampling issues near the poles of extraction spheres sampled equally in standard (θ, ϕ) spherical coordinates.

B. Hydrodynamical initial data

For the simulations considered in this paper, we generate initial data with the Lorene initial data library [105]. We consider three different hydrodynamical configurations. First, an equal mass model with SFHo EOS, corresponding to a binary with Arnott-Deser-Misner (ADM) mass $M_{\text{ADM}} = 2.672M_{\odot}$, initial orbital frequency 331.275 Hz, orbital angular momentum $J = 7.001M_{\odot}/c^2$, initial separation 40.0 km, baryon mass $M_b = 2.967M_{\odot}$ and gravitational mass $2.7004M_{\odot}$. We perform one simulation with the SFHo EOS with mass ratio $q = 1.2$, with ADM mass $M_{\text{ADM}} = 2.675M_{\odot}$, initial orbital frequency 281.407 Hz, orbital angular momentum $J = 7.212M_{\odot}/c^2$, initial separation 45.0 km, total baryon mass $M_b = 2.970M_{\odot}$ and gravitational mass $2.7002M_{\odot}$. We also perform one equal mass simulation with the DD2 EOS, with ADM mass $M_{\text{ADM}} = 2.672M_{\odot}$, initial orbital frequency 331.567 Hz, orbital angular momentum $J = 7.001M_{\odot}/c^2$, initial separation 40.0 km, baryon mass $M_b = 2.941M_{\odot}$ and gravitational mass $2.7001M_{\odot}$. All configurations are initialized with temperature $T = 0.1$ MeV and assuming cold β -equilibrium.

C. Magnetic field configurations

In order to investigate the impact of toroidal fields, which may develop as the result of instabilities in purely poloidal fields as discussed above, we perform a sequence of runs with purely poloidal and purely toroidal magnetic fields, as well as a run with superposed poloidal and toroidal fields, as a model of the configuration that can be expected to develop in an isolated star over many Alfvén times, which we can expect to have developed by the time that a binary system reaches its final few orbits.

In order to construct a purely poloidal magnetic field we initialize the field through taking the curl of a purely toroidal vector potential defined as

$$A^{\phi} = B_0 \max(p - p_{\text{cut}}, 0)^{n_s} \quad (1)$$

$$(A^x, A^y, A^z) = (-yA^{\phi}, xA^{\phi}, 0), \quad (2)$$

with parameters set as detailed in Paper I. In particular B_0 is set such that the magnetic field strength is 5×10^{15} G. This field strength is relatively strong, approaching magnetar strength. The strength of this magnetic field allows us to investigate the strong magnetic fields that develop in the

postmerger phase through the Kelvin-Helmholtz and magnetorotational instabilities (see Paper I) at the finite resolutions we consider, which limit the extent to which amplification through these instabilities can be captured. These choices are consistent with those employed in comparable studies in the literature. We refer to this configuration as POL-UU. To investigate the impact of a toroidal field we use the vector potential specified in Eq. (3), similar to that used in [62,106],

$$\begin{aligned} \mathbf{A} = & B_0 \max(p - p_{\text{cut}}, 0)^{n_s} \\ & \times (x(z^2 - R^2), y(z^2 - R^2), -z(x^2 + y^2 - R^2)). \end{aligned} \quad (3)$$

For the initial data considered in Sec. II B, the stars are initially nonspinning. In long term simulations of magnetized nonspinning stars, there is no process expected to globally break the rotational symmetry, and create a toroidal field with a globally coherent orientation. This is in contrast to a rotating star, whose angular momentum could be expected to break symmetry and give a global orientation to the toroidal field profile. The choice of toroidal initial data given in Eq. (3) is antisymmetric over the z axis, and so has no overall favored orientation within the star itself. The binary system, in contrast, has an initial orbital angular momentum, picking out a global orientation which the toroidal field in the remnant can be expected to adopt. We therefore perform two runs with purely toroidal field orientations to assess the impact of the relative orientation of the toroidal field and the orbital angular momentum. The first run, denoted TOR-PP, has the parameter B_0 positive in both stars, while the configuration TOR-PM sets the parameter B_0 positive in the first star, and negative in the second star.¹

Finally we consider the orientation MIX. For this magnetic field configuration we superpose the POL-UU and TOR-PP configurations, scaling the initial strength parameters B_0 to obtain a relative field strength of 85% poloidal to 15% toroidal, approximately matching the field strengths obtained in single star simulations after the saturation of the initial varicose and kink instabilities after the initial few Alfvén times, such as those seen in [51]. We do not expect this configuration to be an exact match for the field rearrangement at late times in an isolated single star, but to give some insight into the impact of a mixed poloidal and toroidal field in a binary merger.

A generic binary system containing poloidal fields may be expected to contain fields not aligned with the orbital angular momentum. To investigate the impact of moving beyond a symmetric magnetic field configuration we consider modifications of the poloidal configuration also.

¹We denote these configurations as P (plus) and M (minus) rather than U (up) and D (down), since the antisymmetry over the $z = 0$ plane does not define a natural direction of the toroidal field orientation in the inspiral.

We consider a configuration POL-UD, with the same magnetic field configuration as POL-UU, with the orientation of the field in one star flipped with respect to the z axis. The configuration POL-90D has both poloidal magnetic fields rotated to lie in the orbital plane ($z = 0$), initially pointing towards each other along their line of sight. The configuration POL-ASYM takes the configuration of POL-UU and shifts the field by 147m in the positive z direction, thus breaking the symmetry of the configuration. Configurations POL-Q12, STIFF, LOWB and BITANT all have the same magnetic field configuration as POL-UU, but differ due to a mass ratio $q = 1.2$, a DD2 EOS, a low initial field strength $B_0 = 5 \times 10^8$ G and an enforced bitant symmetry, respectively. All configurations other than POL-Q12 are equal mass. Other than BITANT, all runs have no imposed grid symmetry. We also evolve a reference configuration B0 with no magnetic field. In configuration POL-Q12 we emphasize that while the central value of the magnetic field is shared between the two stars, the interior configuration will differ between the two stars, due to the dependence of the vector potential on p . Since the magnetic field strength in a NS is largely determined by rotational properties and the thermal history of the NS, we do not expect a significant difference in the interior maximum strength between the two stars of different masses.

III. RESULTS

A. Gravitational waves

We first present the gravitational waveforms from the performed simulations. We present the gravitational wave strain obtained through the double time integration of the Newman-Penrose scalar Ψ_4 , directly extracted from the numerical simulation on spheres of radius $r_{\text{ex}} = 400M_\odot \approx 588$ km centered around the center of mass of the system. The integration is performed using the fixed frequency integration of [107]. We define the GW strain h , amplitude A , phase ϕ and instantaneous frequency ω , as decomposed onto spherical harmonics of spin weight -2 , and retarded time coordinate u ,

$$\Psi_{4,(\ell,m)} = \ddot{h}_{\ell,m}, \quad (4)$$

$$r_{\text{ex}} h_{\ell,m} = A_{\ell,m} e^{-i\phi_{\ell,m}}, \quad (5)$$

$$\omega_{\ell,m} = \dot{\phi}_{\ell,m}, \quad (6)$$

$$u = t - R - 2M \log \left(\frac{R}{2M} - 1 \right), \quad (7)$$

$$R = r_{\text{ex}} \left(1 + \frac{M}{2r_{\text{ex}}} \right)^2, \quad (8)$$

where M is the gravitational mass.

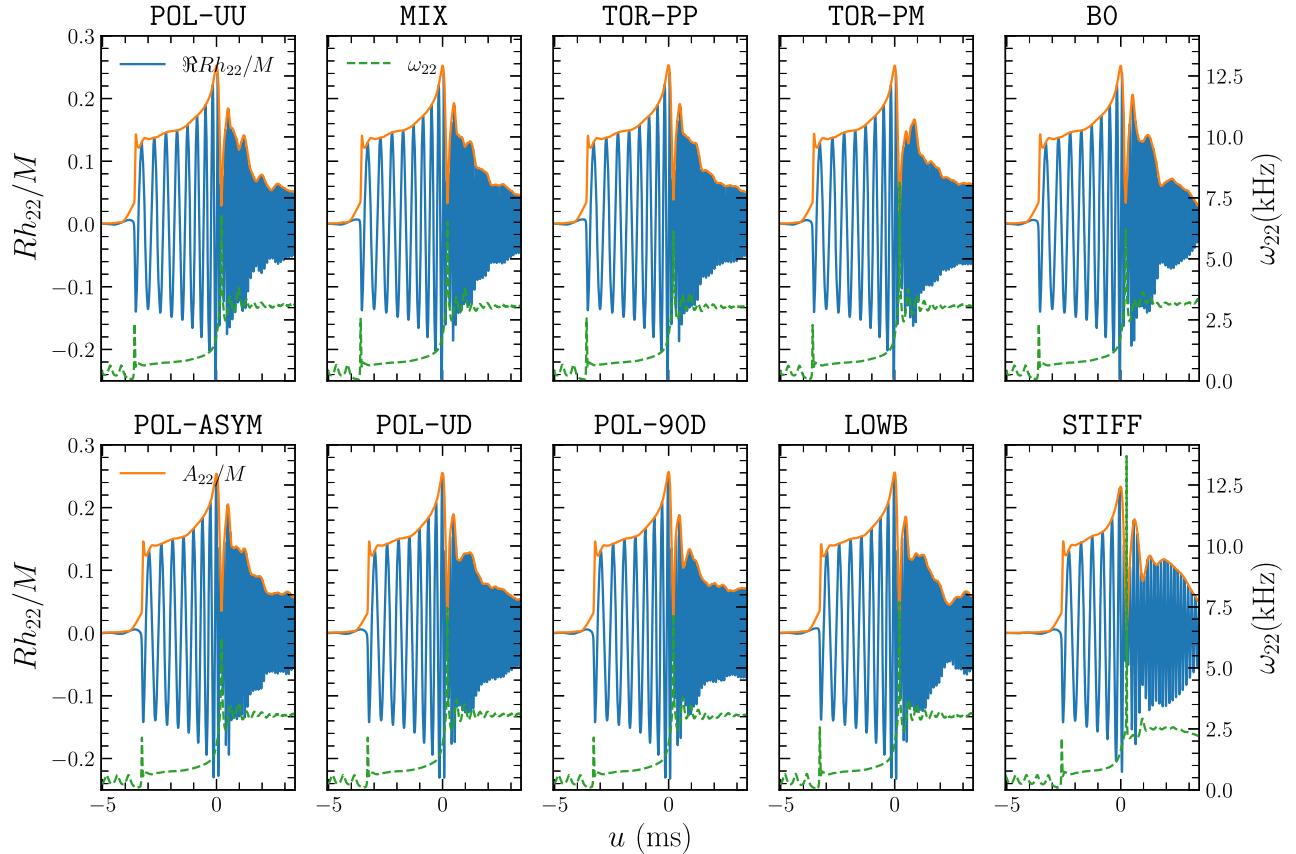


FIG. 1. The real part of the $(\ell, m) = (2, 2)$ mode of the gravitational wave strain h_{22} and its overall amplitude (left axes) and its instantaneous frequency (right axes); extracted at a radius of 588 km.

We demonstrate the GW strains for the equal mass configurations in Fig. 1. Due to the comparatively small energy density in the magnetic field compared to the rest mass energy of the NSs, the inspiral phase is almost identical between the considered simulations, with a close match in the GW signals during this phase. All equal mass simulations complete the same number of orbits (3.5), other than the *STIFF* configuration, which completes half an orbit fewer (3), irrespective of any difference in magnetic field evolution. In the postmerger phase, however, we see some variations in the GW signal for different magnetic field configurations, for different GW modes, as also seen by [62,63]. For the dominant $(\ell, m) = (2, 2)$ mode, we see that the postmerger behavior is broadly similar between all configurations, with the gravitational wave frequency increasing postmerger, as the remnant star oscillates, as the overall amplitude decays. While most configurations show a clear secondary peak in the GW amplitude postmerger, the run *TOR-PM* demonstrates an initial subdominant peak before falling and rising again to a second peak, which is the largest amplitude for this configuration postmerger. We note that this delayed peak is correlated with the delayed amplification of the maximum magnetic field strength in this configuration, see Sec. III B 1. The instantaneous frequency evolution of the postmerger GW in the $(2, 2)$ mode is relatively consistent between configurations with different magnetic field geometries and different magnetic field strengths, and is dominated by the hydrodynamic parameters of the binary. In all SFHo cases, regardless of magnetic field topology, we see a peak of the gravitational wave frequency at merger before the frequency quickly reduces to 3.06 kHz, while for the *STIFF* EOS we see a lower postmerger frequency of 2.36 kHz. These values are consistent with those found using other numerical codes, e.g., [108,109], and in line with expectation. The softer EOS SFHo produces more compact remnants, and thus higher frequency postmerger oscillations of the cores, as opposed to the stiffer DD2 EOS, with less compact remnants. Postmerger, we see subdominant modes more excited due to asymmetries in the system, and the presence of a magnetic field excites higher modes more strongly at merger than the unmagnetized case, due to enhancements of this asymmetry. Immediately after merger the $(2, 1)$ and $(3, 3)$ modes are noticeably excited with amplitudes $\mathcal{O}(1\%)$, the strength of the dominant $(2, 2)$ mode. Notably, in the unmagnetized case, these modes are ~ 3 times smaller than in the magnetized configurations. It has been proposed by [110] that measurements of both the $(2, 1)$ and $(2, 2)$ modes in the postmerger phase may provide a constraint on the energy density gap of a potential QCD phase transition in the merger remnant. This phase transition acts to weaken the one-armed spiral instability in the remnant and consequently the $(2, 1)$ mode. Reference [110] considers evolutions without a magnetic field, and the enhancement of the $(2, 1)$ mode

that we see as a result of the presence of a magnetic field may affect the detectability of this feature, though a further study with EOS incorporating phase transitions, as well as magnetic field evolution, would be required to make a conclusive statement. When comparing to the *LOWB* configuration, we see that, while for the first ~ 0.5 ms postmerger the evolution of the subdominant modes appears similar to the *B0* configuration, these modes are quickly excited, by ~ 1 ms postmerger, to the level of the same modes of the fully magnetized configurations. The *STIFF* configuration also shows a similar excitement of these higher modes at merger to the unmagnetized SFHo case, as it is more resistant overall to the induced oscillations, due to the stiffer EOS.

It is clear that, while the dominant $(2, 2)$ mode frequency is chiefly determined by hydrodynamic quantities such as EOS, mass ratio and total mass, the dynamics of the energy density associated to the magnetic field, subdominant to the total rest mass energy, can have a nontrivial effect on higher modes, due to the induced asymmetry of this energy density. This impact is more visible in these modes, due to their lower overall excitation compared to the $(2, 2)$ mode.

As a consequence the phase evolution of the subdominant GW modes for different magnetic field configurations differs postmerger. As discussed below in Sec. III B, shortly after merger we see a strong dependence of the magnetic field evolution on the initial configuration. The differing distribution of the magnetic field and the associated pressure and tension at this time may have an impact on the density evolution, and consequently, the GW emission in this period immediately postmerger. By comparing the phase evolution of a given GW mode with the reference $(2, 2)$ mode, we can provide a measurable impact of the magnetic field on the GW emission. Specifically we consider the $(2, 1)$ and $(3, 3)$ modes, since they are relatively high in amplitude (a factor of $\sim 50\times$ smaller than A_{22}), and consider their dephasing with respect to the reference $(2, 2)$ mode. At the moment of merger, denoted by the peak in GW strain amplitude, we set the GW phase to zero as a reference, and track the phase difference accumulated between the $(2, 2)$ mode and these subdominant modes, up until the time at which the GW amplitude drops to half of its postmerger maximum. In Fig. 2 we show a scatter plot of these dephasings $\Delta\phi_{\ell,m} = \phi_{\ell,m} - \phi_{22}$.

We see that, compared to the unmagnetized case *B0*, the presence of a magnetic field may induce a dephasing between the $(2, 1)$ [$(3, 3)$] and $(2, 2)$ modes of up to 22.5 (8.2) radians by the time the GW signal has decayed to half of its postmerger peak, maximized by the *POL-ASYM* configuration, with the largest initial magnetic field asymmetry, while over all different magnetic field configurations an overall difference in the dephasing of up to 30.5 (14.6) radians can be reached. There is no clear hierarchy

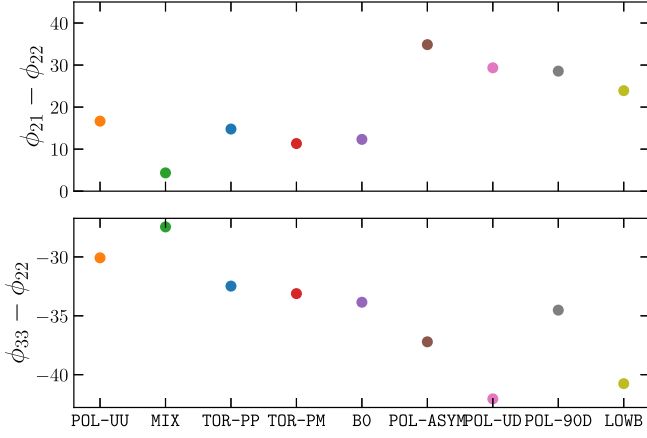


FIG. 2. Upper (Lower) panel: The dephasing between the (2, 2) mode and the (2, 1) ((3, 3)) mode $\Delta\phi_{21}$ ($\Delta\phi_{33}$).

visible between the B0, LOWB and remaining strongly magnetized configurations.

We compare the dephasing effects seen in our reference SR configurations with two reference runs at a lower resolution for the POL-UU and B0 configurations to estimate the error in these effects. For these low resolution runs we find a difference in $\Delta\phi_{21}$ between the POL-UU and B0 configurations of 3.0 radians, and 5.9 radians for $\Delta\phi_{33}$. The corresponding dephasings for our SR configurations are 4.3 and 3.8 radians, respectively. These provide a relative error estimate of 30% and 55%, respectively, for the dephasing in the (2, 1) and (3, 3) modes arising from finite resolution effects. We, however, see a much larger error arising from the choice of reconstruction method. In Fig. 2 we see that the dephasings of the configurations performed with the WENOZ reconstruction, and with the PPM reconstruction, form two distinct groups. We perform a direct comparison of the same configuration with these two reconstruction methods in Appendix, and quantify the error in $\Delta\phi_{21}$ ($\Delta\phi_{33}$) as 70% (38%), respectively. However, the absolute size of variation in $\Delta\phi$ arising from the choice of reconstruction is, for the (3, 3) mode, larger than the range of dephasings seen above for different magnetic field configurations in this mode. The size of the variation in dephasings for different magnetic field choices for a given reconstruction choice is approximately constant between the two reconstruction choices considered here, however, it is clear that the dephasings seen are highly dependent on the reconstruction routine used in the numerical evolution.

B. Magnetic field evolution

1. Magnetic field amplification

We now consider the global amplification of the magnetic field, and direct the reader to Paper I for a discussion of the magnetic field in the postmerger remnant and disk. At the initial moment of merger we see the amplification of the magnetic field through the onset of the KHI. Here a

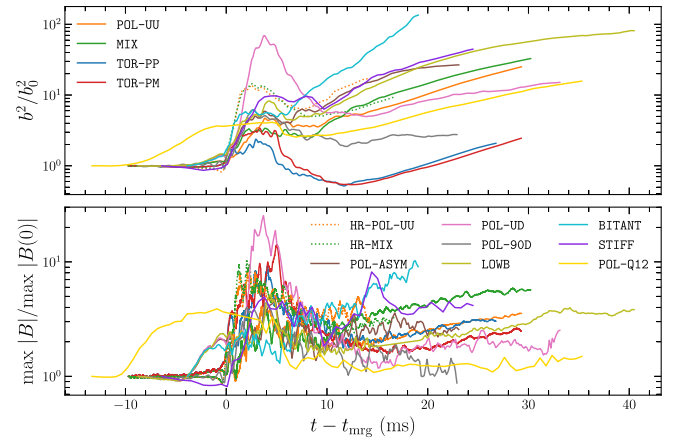


FIG. 3. Upper panel: the amplification of the magnetic field energy. Lower panel: the amplification of the maximum magnetic field strength.

shearing layer forms between the two stars which breaks down, forming small vortices in the fluid flow. These vortices wind up the magnetic field lines, pinned to the fluid flow within the ideal MHD approximation, leading to an amplification of the magnetic field strength. From high resolution simulations [5] it is known that this instability can amplify the magnetic field energies by in excess of 6 orders of magnitude, with the limiting factor being the resolution with which the simulations are performed. Finite resolution effects restrict the size of the vortices that can be captured, and so restrict the scale at which the amplification due to the KHI cannot be resolved beneath.

We measure the dependence of this instability on our initial field geometry through, first, the total energy amplification and, secondly, the amplification of the pointwise maximum magnetic field strength, shown in Fig. 3. In the upper panel of Fig. 3 we see the amplification of the magnetic field energy as a proportion of the initial energy. First let us focus on the amplification arising from the KHI, centered around the moment of merger $t = 0$. The fiducial run POL-UU amplifies by a factor ~ 5 by 4 ms postmerger. In contrast, the configurations initialized with toroidal initial data see a suppressed amplification of the magnetic field. In the configurations TOR-PP and TOR-PM the magnetic field is maximized not in the equatorial plane, but at nonzero values of the z coordinate. When the stars first merge, and the shearing layer associated to the KHI first forms, it is initially formed in the equatorial plane, as this is the first point of contact between the stars. For the toroidally dominated configurations, the magnetic field is comparatively weaker at this location, and so the amplification due to the KHI is relatively weaker, explaining the difference in behavior for these runs.

We note that the higher resolution simulations of the POL-UU and MIX configurations demonstrate, as expected, a considerably larger amplification, amplifying twice as much as their lower resolution counterparts at the moment

of merger, but demonstrating qualitatively similar behavior at the moment of merger. This is consistent with behavior seen in resolution studies of the Kelvin-Helmholtz instability in BNS mergers, for instance [5], who find the growth rate of the instability proportional to the inverse of the discretization length. Since the TOR-PP configuration demonstrates very little amplification, the addition of such a component to a poloidal component has very little effect, with the MIX configuration and POL-UU configuration behaving very similarly.

Breaking the symmetry of the poloidal magnetic field configuration, however, has a considerable effect. When the magnetic fields are totally antialigned, and poloidal, for configuration POL-UD, we see the largest amplification, by a factor of ~ 100 , with the other antialigned configuration, TOR-PM, also amplifying considerably more than the aligned TOR-PP configuration, though the effect here is much less stark due to the overall lower amplification due to the toroidal field.

Configurations with antialigned magnetic fields may be considered as candidates for enhanced reconnection of magnetic field lines due to the relative orientation of the fields. As a proxy for this we calculate the current $\mathbf{J} = \nabla \times \mathbf{B}$, and find that the hierarchy of the current matches that of the magnetic field amplification. In local simulations of the KHI it has been found that the growth rate of the instability is enhanced for a magnetic field that switches direction over the shearing interface, as in our antialigned configurations POL-UD and TOR-PM, compared with a reference configuration with a uniform field over the interface [111]. These simulations are performed for resistive Newtonian MHD, and here we find the same behavior in the case of ideal general relativistic (GR) MHD. It is suggested in [111] that this enhancement arises from the increased reconnection, which forms field lines with increased magnetic tension, contributing to the development of the KHI. In our simulations, we also find that the hierarchy of the magnetic tension in the star matches that of the magnetic field and current. Since our simulations are in ideal MHD we do not include reconnection effects, but we may expect numerical reconnection to occur due to artificial dissipation arising from our numerical approach. Any amplification of the magnetic field arising from this effect may be expected to compete with a loss of magnetic field energy in reconnection events, which should convert magnetic field energy into kinetic energy; though clearly here we see that the field amplification is dominant in these cases with antialigned fields.

Immediately after the initial amplification associated to the KHI, we see an intermediate phase before the onset of a second phase of amplification, driven by winding, which lasts between ~ 5 – 10 ms postmerger. For most configurations the magnetic field energy remains approximately constant during this phase, with the following exceptions. All toroidally dominated runs see a considerable

dissipation of magnetic field energy between ~ 4 – 12 ms postmerger, as does POL-UD dropping down to a similar level to configuration POL-UU by the end of this intermediate phase. This behavior has previously been observed [50] and can arise from the small scale structures formed during the KHI, associated to the amplified magnetic field, dropping below the grid resolution, with numerical dissipation effectively dissipating some of the amplified field.

Beginning between 9 and 14 ms postmerger, we see the onset of a winding driven amplification phase. During this phase we expect to see a linear growth in toroidal magnetic field strength, associated to the overall angular momentum of the remnant winding up a toroidal field. We perform fits to the postmerger magnetic field strength and find that the expected linear behavior is recovered for all configurations considered. We find that toroidally dominated configurations show a slower growth rate here, in line with their behavior in the KHI phase, since the radial field that has been developed by this point is smaller, though in Paper I we found similar rotation rates between all models in the remnant star.

We note that configurations STIFF and POL-Q12, with the same field as POL-UU, demonstrate qualitatively similar behavior to that run, suggesting that the amplification is largely insensitive to the hydrodynamic configuration. As observed in Paper I, we see that, by the end of the BITANT configuration, it has amplified by a factor of 100 times more than the least amplified cases, the TOR-PP and TOR-PM configurations. In Paper I we have suggested that the bitant symmetry enhances amplification by maintaining poloidal fields at the symmetry plane, enhancing the winding mechanism for field amplification. We note that in [60] a direct comparison of a configuration with and without bitant symmetry is made for a different magnetic field configuration extending outside of the star, and in contrast to our simulations, similar magnetic field amplifications are seen in both cases.

In the lower panel of Fig. 3 we see a consistent hierarchy also with the pointwise maximum of the magnetic field. We note here that configuration TOR-PM shows a delayed amplification to its maximum, correlated with the delayed secondary peak in the gravitational wave amplitude seen.

2. Toroidal field evolution

In Fig. 4 we track the development of the toroidal field inside a single star during the inspiral and merger, focusing on the MIX configuration, with the POL-UU and toroidally dominated configurations for comparison. The MIX configuration is constructed to approximate the end state of the dynamical instabilities that purely poloidal configurations such as POL-UU are known to suffer from, which generate toroidal fields. We note that Paper I contains a full discussion of the toroidal field within the remnant, and so here we focus on the relevance to the choice of initial data during the inspiral. We plot $b_\phi b^\phi$

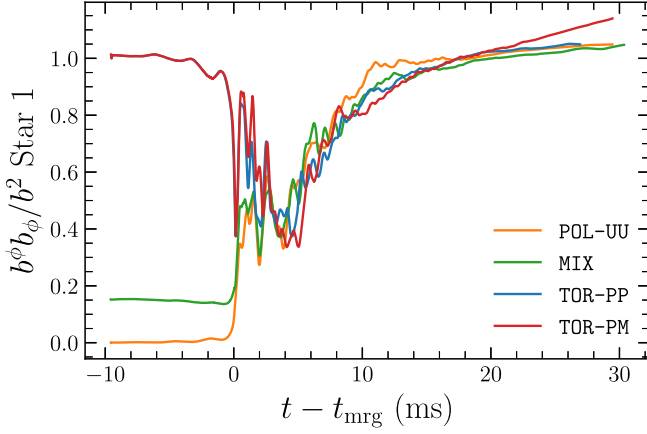


FIG. 4. The toroidal magnetic energy as a proportion of the total magnetic energy of a single star.

normalized by the norm of the four-vector b^2 , noting that this quantity can exceed 1. We see that during the inspiral the toroidal component of the magnetic field within the stars is largely unchanged, and up until ~ 2 ms before merger the field is the same as that set in the initial data. This suggests to us that the relatively short inspiral considered in this work, consisting of 3 orbits and lasting 10 ms is insufficiently long to observe the effects of the dynamically developing magnetic field instabilities expected in configuration POL-UU. In [51] such instabilities are observed to saturate on the order of 25 ms, albeit for a stronger magnetic field strength, with a correspondingly faster Alfvén time. For the magnetic field strengths considered in this study, one might expect these instabilities to develop in the premerger stars over the course of $\mathcal{O}(100 \text{ ms})$, driving the POL-UU configuration towards a configuration similar to the MIX configuration. At the moment of merger the poloidally dominated configurations begin to develop a toroidal field, while the toroidal configurations begin to lose their toroidal field, as the localized fields given by the initial configurations are disrupted, and replaced with the gradually forming global toroidal field. This is driven by the rotating remnant, with the field growing at an approximately universal rate over these configurations, due to the common rotation rate between models (Paper I).

3. Poynting flux

After the merger, energy is radiated through the electromagnetic field through the Poynting flux. We calculate the luminosity associated to the radial Poynting flux, defined as

$$L_{\text{Poynting}} = \frac{dE}{dt} = \int_{\Omega} -T_{t\text{EM}}^r r^2 d\Omega = \int_{\Omega} (b^r b_t - b^2 u^r u_t) r^2 d\Omega, \quad (9)$$

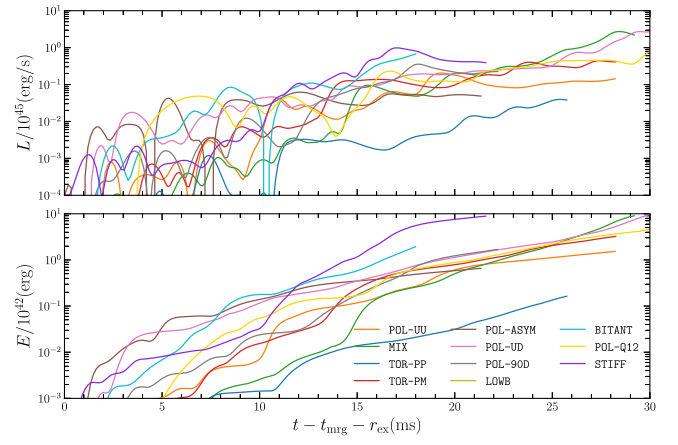


FIG. 5. Upper panel: the instantaneous luminosity of the radial Poynting flux in units of $10^{45} \text{ erg s}^{-1}$. Lower panel: the cumulative energy in ergs radiated through the radial Poynting flux up to time t . The time shown is the time in milliseconds postmerger, less the extraction radius.

which we calculate on an extraction sphere Ω with a radius of 295 km centered at the origin. We show the luminosity of the Poynting flux, and its time integral, the radiated energy, in Fig. 5.

The evolution of the Poynting luminosity is broadly similar across different magnetic field configurations. We find that the strength of the flux is largely correlated with the strength of the late time magnetic field in the remnant, developed through winding, expected as the source of the power of the late time Poynting flux, with the BITANT, STIFF and MIX configurations, in particular, showing the strongest fluxes.

Measured 15 ms postmerger these configurations can reach luminosities a factor $\sim 10^3$ larger than that of the weakest Poynting flux, while, amongst the remaining configurations, the luminosity may vary by factors $\mathcal{O}(10^2)$. Configuration POL-ASYM, which has a large amplification in the winding phase, has a relatively low Poynting luminosity by the end of the simulation. We note that, at this time, its amplification due to winding has slowed and that its peak in Poynting luminosity, ~ 10 ms postmerger, is coincident with its largest magnetic field amplification in the winding phase. This supports the claim that the winding driven magnetic field is determining the Poynting luminosity.

The LOWB run has too low a Poynting flux to be visible in Fig. 5, though we note that, measured at the same time postmerger, its luminosity is consistently a factor $\sim 10^{12}$ lower than that of the POL-UU configuration. This factor is broadly consistent with the expected scaling of the Poynting flux with B_0^2 , the initial strength of the magnetic field [112,113].

Comparing to much longer lasting and higher resolution simulations such as those presented in [114], we can see

that the Poynting fluxes in our simulations are consistent with these simulations over the time span we consider postmerger, approaching an approximately steady phase after tens of ms. At later times, closer to 100 ms postmerger, [114] sees the formation of a far stronger Poynting flux. Based on the strength of the initial magnetic field in our simulations, we may expect longer lived extensions of our simulations to reach luminosities on the order of 10^{51} erg/s [113], and see an accompanying effect driving the ejection of matter from the merger remnant.

C. Symmetry

High resolution NR simulations of BNS systems, evolving a range of physics, employing tabulated EOS, can be incredibly expensive. In order to reduce computational cost, symmetry assumptions can be made to reduce the extent of the computational domain. A natural choice for the binary problem is to assume bitant symmetry, assuming a symmetry over the orbital plane of the binary (here the $z = 0$ plane), evolving only the upper half plane of the computational domain. The evolution is then performed using appropriate boundary conditions over the $z = 0$ plane. Scalar quantities such as the density are set such that $\rho(z = 0^+) = \rho(z = 0^-)$, the velocity is set such that $(\tilde{u}^x, \tilde{u}^y, \tilde{u}^z)(z = 0^+) = (\tilde{u}^x, \tilde{u}^y, -\tilde{u}^z)(z = 0^-)$ and, for consistency with the standard assumption of an internal poloidal field in the star, the magnetic field boundary conditions are set such that $(B^x, B^y, B^z)(z = 0^+) = (-B^x, -B^y, B^z)(z = 0^-)$. We perform all of the runs considered in this paper assuming no symmetry at all, except for BITANT, which is initialized with the poloidal field configuration, and employs the above mentioned symmetry over the $z = 0$ plane. We analyze the extent to which our runs with no enforced symmetry retain the symmetry properties of the initial data, and assess to what extent an enforced bitant symmetry can give a false impression of the overall magnetic field structure which develops during the simulation. For these purposes we define the following diagnostic quantities for the symmetric and antisymmetric parts of the magnetic field:

$$B_{z>0}^i = \frac{\int_{z>0} B^i \rho dV}{\int_{z>0} \rho dV}, \quad (10)$$

$$B_{z<0}^i = \frac{\int_{z<0} B^i \rho dV}{\int_{z<0} \rho dV}, \quad (11)$$

$$B_{\text{sym}}^i = \frac{1}{2} (B_{z>0}^x - B_{z<0}^x, B_{z>0}^y - B_{z<0}^y, B_{z>0}^z + B_{z<0}^z) \quad (12)$$

$$B_{\text{asym}}^i = \frac{1}{2} (B_{z>0}^x + B_{z<0}^x, B_{z>0}^y + B_{z<0}^y, B_{z>0}^z - B_{z<0}^z) \quad (13)$$

$$B_{\text{sym}} = \sqrt{B_{\text{sym}}^i B_{i\text{sym}}} \quad (14)$$

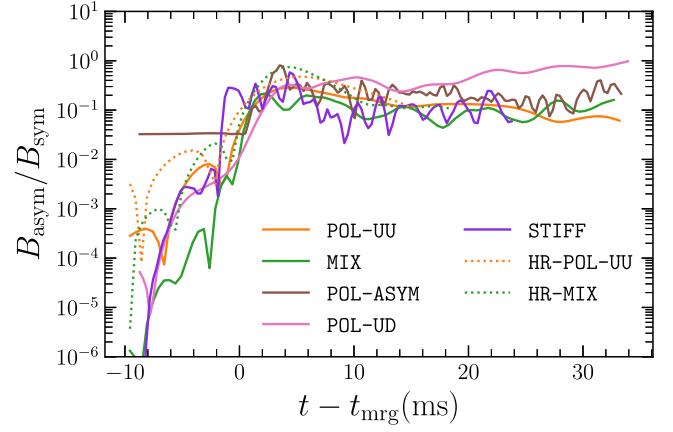


FIG. 6. The ratio between the antisymmetric and symmetric parts of the magnetic field integrated over the spatial domain [Eqs. (14)–(15)].

$$B_{\text{asym}} = \sqrt{B_{\text{asym}}^i B_{i\text{asym}}}. \quad (15)$$

In order to avoid polluting our results with fields supported by spurious low density ejecta, we only integrate over matter with density over 10^{13} g/cm³.

In Fig. 6 we present the time evolution of the symmetry of the magnetic field. We comment that the toroidal and POL-90D configurations are not adapted to the symmetry considered, and so we do not include them here, and that by construction, the POL-ASYM configuration begins with an asymmetry. Clearly we see at early times the magnetic field is largely symmetric for the configurations adapted to the bitant symmetry, with a relative initial violation of symmetry on the order of one part in 10^3 if not smaller. We see that asymmetry grows throughout the inspiral phase, and that by the point of merger, the antisymmetric component of the field has grown to several percent the size of the symmetric field. This behavior is largely independent of the specific initial configuration, with the rate of increase of antisymmetry throughout the inspiral approximately uniform for different configurations. For the POL-ASYM configuration, which begins at a nontrivial value of the asymmetry by construction, once sufficient time has passed during the inspiral for the initially symmetric configurations to grow an asymmetry of a similar size to that of the initial data of POL-ASYM, the asymmetry evolves along the same trajectory as the initially symmetric runs. This suggests that the growth of the asymmetry is occurring universally across all runs, but is only visible in the POL-ASYM configuration once it exceeds the initial asymmetry. We note that for the runs not presented here that are not adapted to the bitant symmetry, that while the asymmetry may be large in the inspiral due to the construction of the initial data, the magnitude in the postmerger phase is consistent with the other runs. This is in agreement with the findings of Paper I, which suggests that the toroidal

field in the postmerger remnant is largely consistent between configurations.

In the immediate postmerger phase the asymmetry grows to order 10% of the symmetric field, and remains approximately at this level, similar for all configurations considered. At these values, approximating the true configuration with a bitant symmetry can be expected to be a poor model for the magnetic field structure. The most extreme configuration, POL-UD, grows its asymmetric field up to order unity in size compared to the symmetric field by the end of the simulations considered, at which point a bitant symmetry would not be capturing the true magnetic field at all. In the benchmark POL-UU configuration, we see that the asymmetry begins to decay slightly in the postmerger phase, below 10%. This standard configuration appears to match the bitant symmetry the closest at late times, though the value of the asymmetric field is still much higher than one would hope for if a bitant symmetry run were expected to capture the full magnetic field dynamics.

Despite the fact that the initial data is symmetric, and the evolution system introduces no explicit breaking of this symmetry, we have seen a nontrivial asymmetry grow throughout our evolution, that appears to be approximately uniform in its growth rate and final saturation value, independent of the specific initial condition. This behavior is consistent with the onset of a spontaneous symmetry breaking bifurcation, with the solution to the Euler equations undergoing a transition to a lower symmetry than that of the initial conditions (see, for instance, [115] for a review of such phenomena). The growth rate of the instability should then be driven by the growth rate of modes close to the bifurcation, independent of the initial data itself. The imposition of bitant symmetry prevents this symmetry breaking from occurring, and so associated phenomena will be neglected in such simulations. In particular, in Paper I we discuss a suppressed formation of turbulent features in the BITANT configuration. Here we see that the formation of such features in the magnetic field present in full grid symmetry, but not in the BITANT symmetry, can lead to a nontrivial breaking of the symmetry in the magnetic field.

D. Ejecta

The binary merger process ejects neutron rich matter from the postmerger remnant, which will decay to stable isotopes, giving off a long-lived electromagnetic signal, the kilonova [18–20]. This process is believed to be the origin of a large proportion of the heavy elements found in the Universe [16,17]. The presence of magnetic fields may create additional pressures in the postmerger remnant, which can affect the amount of material ejected from the merger, potentially enhancing the material undergoing r-process nucleosynthesis, and contributing to the kilonova. In addition, the nature of the magnetic field itself within the ejecta may influence the weak processes that are energetically favorable to occur during the decay of the neutron rich

matter. The magnetic field configuration may also affect the long term evolution of the ejected matter, trapping charged particles close to the remnant, or beaming them quickly away [72].

1. Overview

We first discuss the overall evolution of the ejected material. In Figs. 7 and 8 we show the evolution of the ejected material at large distances at two time slices, ~ 6.5 ms postmerger, and the final time in the simulation. We note that the final time of each run is not the same, we refer the reader to Table 1 in Paper I for the duration of each simulation. Due to the different reconstruction methods used (see Appendix), the point 6.5 ms postmerger is not exactly identical between runs, however, the state of the evolution is closely comparable. In the upper panel of each plot we demonstrate the specific entropy per baryon, and in the lower panel demonstrate $\beta^{-1} = b^2/2p$, the ratio between the magnetic and fluid pressures. In white, we demonstrate contours of density, while the thick purple contour demonstrates the temperature contour at 4 GK. This is the threshold at which ejected material has cooled sufficiently to drop out of nuclear statistical equilibrium.

We focus initially on the evolution in the x - z plane in Fig. 7. Across all configurations, shortly postmerger, we see ejecta expanding outwards, with high entropy shock driven material emitted in the polar regions, and lower entropy tidally ejected material emitted in the equatorial plane; as conventionally seen in such mergers, e.g., [116]. On the diagonal of the plot, we see a region of lower temperature and lower entropy material that has dropped below 4 GK, in all configurations, which lies in between these two regions and ejection channels. These low temperature regions correspond also to low pressure regions, and so β^{-1} is also larger in these diagonal regions, with the largest β^{-1} found for the strongest amplifying BITANT and POL-UD configurations. The magnetic field strength itself in this region is consistent with surrounding regions, and the increased β^{-1} only arises due to the reduced p . These low pressure regions are also present in the B0 configuration, though less strongly than in the magnetized cases, and so appear not to be totally magnetically driven. From the magnetic field lines we see that the magnetic field is unstructured in the ejected material with no overall global configuration having formed.

By the final time we see that, a funnel of higher entropy shocked material has formed in the polar region. For the configurations MIX and POL-UD, which amplify fastest and most strongly in the winding phase, respectively, we see high β^{-1} material emitted vertically outwards from the two lobes of the remnant structure in the center, the formation of which we discuss below. The overall field structure at the end of our simulations is still not coherent with no overall structure visible.

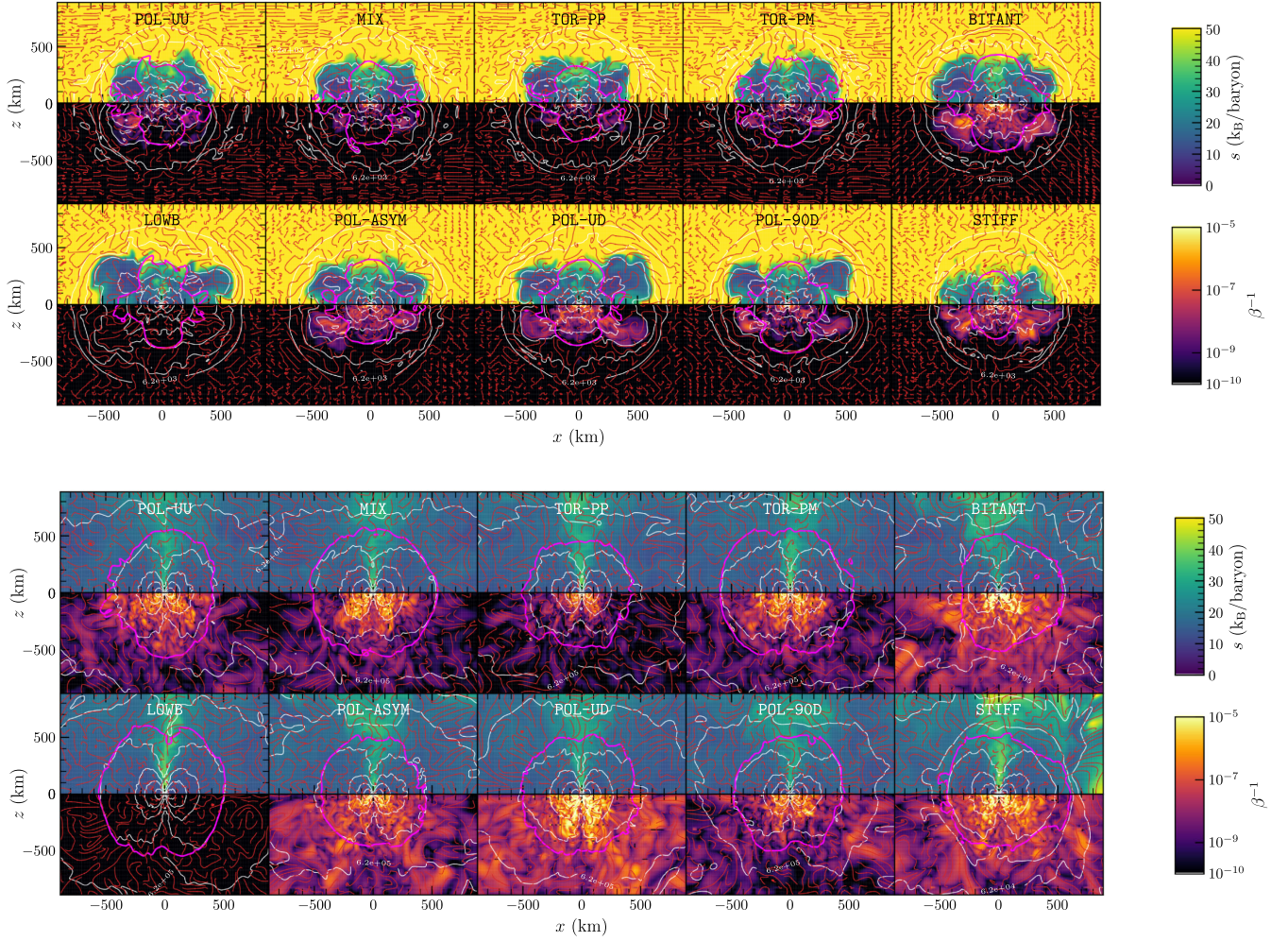


FIG. 7. Two-dimensional (2D) slices in x - z plane demonstrating entropy per baryon in the upper half plane and β^{-1} in the lower half plane. White contours are constant density contours, labeled in units of g/cm^3 . The outermost contour is labeled, all subsequent contours increase by 1 order of magnitude moving inwards towards the remnant. The thick purple contour shows a temperature of 4 GK. Magnetic field streamlines are denoted in red. Upper plot: 6.5 ms postmerger. Lower plot: final time in simulation.

We see similar behavior in the x - y plane in Fig. 8, noting that higher β^{-1} material is ejected further in the equatorial plane in the STIFF configuration compared to the SFHo runs, and that at the final time higher entropy material has been ejected to large radii for the STIFF configuration. This supports the idea that stiffer EOS may drive ejecta by stronger tidal effects [117]. The configurations with the largest amplification in the KHI phase, POL-UD and BITANT, along with POL-90D, have ejected the most high β^{-1} material to a large radius. The field itself is also unstructured in the x - y plane.

It has been suggested [69–71] that strong magnetic fields in the site of r -process nucleosynthesis can modify β -decay rates and neutron decay rates, affecting the relevant Urca rates, due to the quantization of Landau levels. Such effects require relatively strong magnetic fields to be present, at least in excess of 10^{14} G, and for strong effects, strengths of $\sim 10^{16}$ G. We locate the

magnetic field strength isocontour at 10^{14} G in Figs. 7 and 8 (not pictured). This contour lies very close to the remnant, and is ~ 10 km in radius. Clearly the magnetic field strengths close to the remnant exceed these values and so, when incorporating neutrinos into such an evolution, the effects of magnetic fields on the Urca rates in these regions must be incorporated. The field strengths extending outside of the remnant region, however, are insufficiently strong to have such an impact on these rates, as the material expands. For further detail on the field in the remnant, we refer the reader to Paper I.

We denote in Table II the values of key diagnostic quantities we find in the ejecta in all of our simulations, which we make reference to in the following analysis.

2. Hydrodynamical evolution of ejecta

We now discuss the composition of the ejected material, and the impact of the magnetic field configuration on it. We

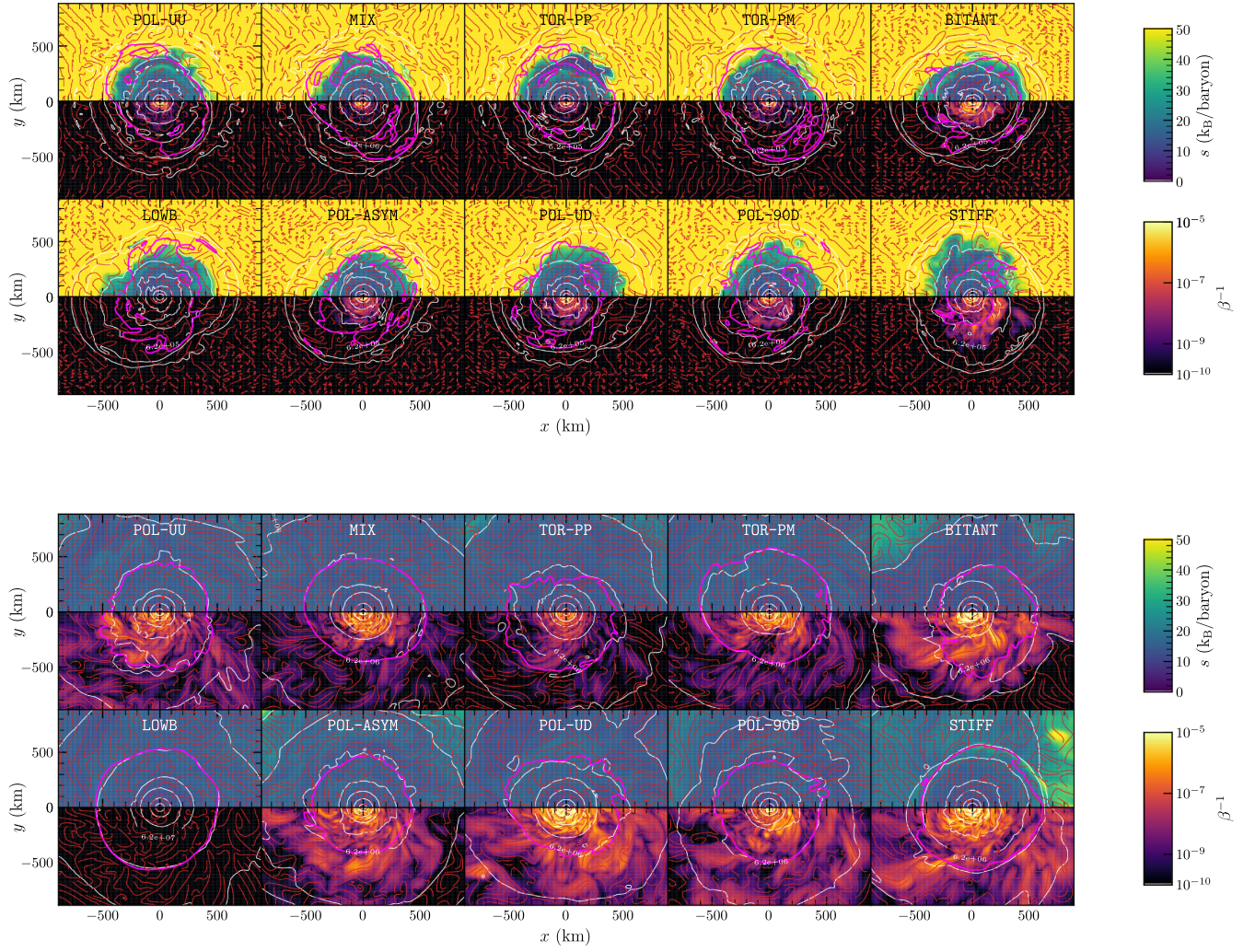


FIG. 8. 2D slices in x - y plane. Upper plot: 6.5 ms postmerger. Lower plot: final time in simulation. Quantities plotted as in Fig. 7.

TABLE II. Summary of ejecta properties for all runs. Columns show final ejected mass, the average value of $\rho, s, v, \theta, B, B^r/|B|, B^\phi/|B|$, the final value of fitting parameters α_1, α_2 defined in Eq. (17), and $\ell_{\max}$ and L_{corr} , the maximum multipole moment and correlation length for the mass weighted toroidal field in the ejecta, discussed in Sec. III D 4. Upper and lower limits denote the interval within which 90% of the ejected mass is located. The LOWB magnetic field strength is amplified by a factor of 10^6 .

Run name	$M_{\text{ej}}(M_{\odot})$	$\langle \rho \rangle$ (10^8g/cm^3)	$\langle s \rangle (k_B/\text{b.})$	$\langle v(c) \rangle$	$\langle \theta \rangle$	$\langle B(10^9 \text{G}) \rangle$	$\langle B^r/ B \rangle$	$\langle B^\phi/ B \rangle$	α_1	α_2	$\ell_{\max}$	L_{corr} (km)
POL-UU	0.027	$0.95^{+0.99}_{-0.81}$	$16.46^{+7.01}_{-7.28}$	$0.18^{+0.15}_{-0.09}$	$1.57^{+1.03}_{-1.08}$	$1.51^{+2.66}_{-1.45}$	$0.00^{+0.81}_{-0.86}$	$0.01^{+0.93}_{-0.99}$	0.94	0.60	3	98.3
MIX	0.027	$0.88^{+0.52}_{-0.74}$	$16.62^{+6.85}_{-7.44}$	$0.19^{+0.17}_{-0.10}$	$1.59^{+1.13}_{-1.18}$	$2.14^{+2.03}_{-2.11}$	$-0.00^{+0.82}_{-0.89}$	$0.03^{+0.91}_{-1.00}$	0.84	0.41	1	295.0
TOR-PP	0.022	$0.94^{+1.01}_{-0.84}$	$16.28^{+7.19}_{-7.10}$	$0.19^{+0.15}_{-0.10}$	$1.59^{+1.00}_{-1.05}$	$0.60^{+0.63}_{-0.58}$	$0.00^{+0.82}_{-0.86}$	$-0.03^{+0.97}_{-0.95}$	1.00	0.69	6	49.2
TOR-PM	0.033	$1.06^{+0.89}_{-0.92}$	$16.32^{+7.15}_{-7.13}$	$0.19^{+0.13}_{-0.09}$	$1.57^{+1.02}_{-1.09}$	$1.23^{+2.94}_{-1.20}$	$0.00^{+0.85}_{-0.90}$	$-0.02^{+0.91}_{-0.96}$	0.92	0.58	17	17.4
BO	0.026	$0.95^{+0.99}_{-0.85}$	$16.38^{+7.09}_{-7.20}$	$0.18^{+0.19}_{-0.09}$	$1.58^{+1.08}_{-1.16}$	...	...	...	...	...	...	...
BITANT	0.013	$1.06^{+0.89}_{-0.92}$	$15.67^{+7.80}_{-6.48}$	$0.22^{+0.12}_{-0.11}$	$1.54^{+1.13}_{-0.99}$	$4.41^{+9.74}_{-4.21}$	$0.02^{+0.84}_{-0.92}$	$0.04^{+0.90}_{-0.98}$	0.65	0.19	7	42.1
LOWB	0.023	$0.80^{+0.60}_{-0.70}$	$16.57^{+8.94}_{-7.39}$	$0.20^{+0.18}_{-0.11}$	$1.55^{+1.18}_{-1.20}$	$1.57^{+1.72}_{-1.53}$	$-0.00^{+0.86}_{-0.89}$	$-0.02^{+0.96}_{-0.96}$	...	...	3	98.3
POL-UD	0.023	$0.74^{+0.66}_{-0.64}$	$16.83^{+8.68}_{-7.65}$	$0.20^{+0.18}_{-0.11}$	$1.60^{+1.13}_{-1.25}$	$5.98^{+8.17}_{-5.87}$	$-0.01^{+0.86}_{-0.89}$	$-0.02^{+0.96}_{-0.96}$	0.75	0.35	5	59.0
POL-90D	0.019	$0.75^{+0.65}_{-0.65}$	$16.52^{+8.99}_{-7.34}$	$0.21^{+0.17}_{-0.12}$	$1.57^{+1.02}_{-1.09}$	$1.65^{+2.52}_{-1.59}$	$-0.01^{+0.86}_{-0.89}$	$0.02^{+0.92}_{-1.00}$	0.92	0.73	4	73.8
POL-ASYM	0.019	$0.78^{+0.62}_{-0.68}$	$16.23^{+7.24}_{-9.08}$	$0.20^{+0.16}_{-0.09}$	$1.59^{+1.01}_{-1.11}$	$2.96^{+4.72}_{-2.85}$	$0.01^{+0.85}_{-0.90}$	$0.03^{+0.91}_{-0.97}$	0.82	0.54	2	147.5
STIFF	0.020	$0.68^{+0.72}_{-0.64}$	$19.18^{+10.41}_{-5.91}$	$0.16^{+0.15}_{-0.07}$	$1.55^{+1.11}_{-1.13}$	$5.51^{+8.64}_{-5.31}$	$-0.01^{+0.78}_{-0.85}$	$0.03^{+0.91}_{-1.01}$	0.82	0.53	17	17.4
POL-Q12	0.010	$0.83^{+6.42}_{-0.78}$	$12.75^{+86.23}_{-7.65}$	$0.22^{+0.28}_{-0.09}$	$1.49^{+1.62}_{-0.88}$	$7.13^{+81.32}_{-6.77}$	$0.01^{+0.97}_{-0.95}$	$-0.09^{+1.07}_{-0.89}$	0.54	0.76	19	15.5

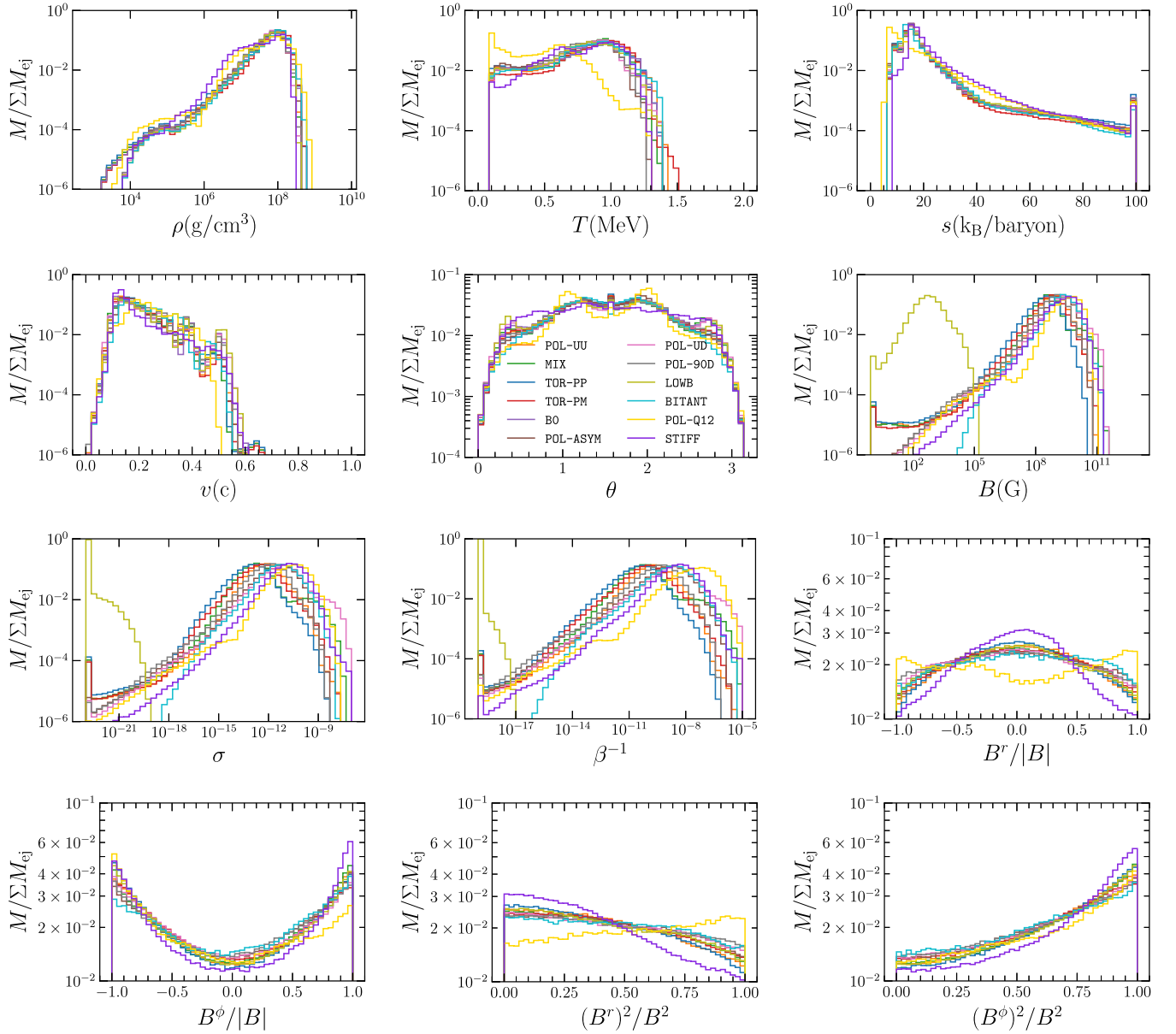


FIG. 9. Histograms binned over various quantities demonstrating the cumulative proportion of the total ejected mass up to the end of the simulation, in those bins. The quantities plotted are from left to right, then top to bottom: density, temperature, specific entropy per baryon, velocity, colatitudinal angle, total magnetic field strength, magnetization, β^{-1} , ratio of radial magnetic field strength to total magnetic field strength, ratio of toroidal magnetic field strength to total magnetic field strength, ratio of absolute radial magnetic field strength to total magnetic field strength. Colors demonstrate different magnetic field configurations.

first demonstrate the mass weighted composition of the ejecta at the end of our simulations in the histograms in Fig. 9. To complement this discussion with a measure of the time evolution of the ejecta properties, we demonstrate also the average value of various diagnostic quantities within the ejecta over time in Fig. 10. In this discussion we classify as ejecta material that is unbound according to the Bernoulli criterion, that is, material as extracted on a sphere of radius 295 km, for which $hu_t < -1$. Here u_t is the timelike covariant component of the fluid 4-velocity,

while $h = 1 + \epsilon + \frac{p}{\rho}$ is the fluid enthalpy. ϵ is the fluid specific internal energy.

In the simulations performed we have neglected the effect of neutrinos, with the electron fraction remaining constant. We may expect that including these effects may affect the total amount of ejected material, and its thermodynamical properties, such as the entropy and temperature [116]. We nevertheless discuss such thermodynamical quantities here to flag the impact of magnetic field topology on the ejected material.

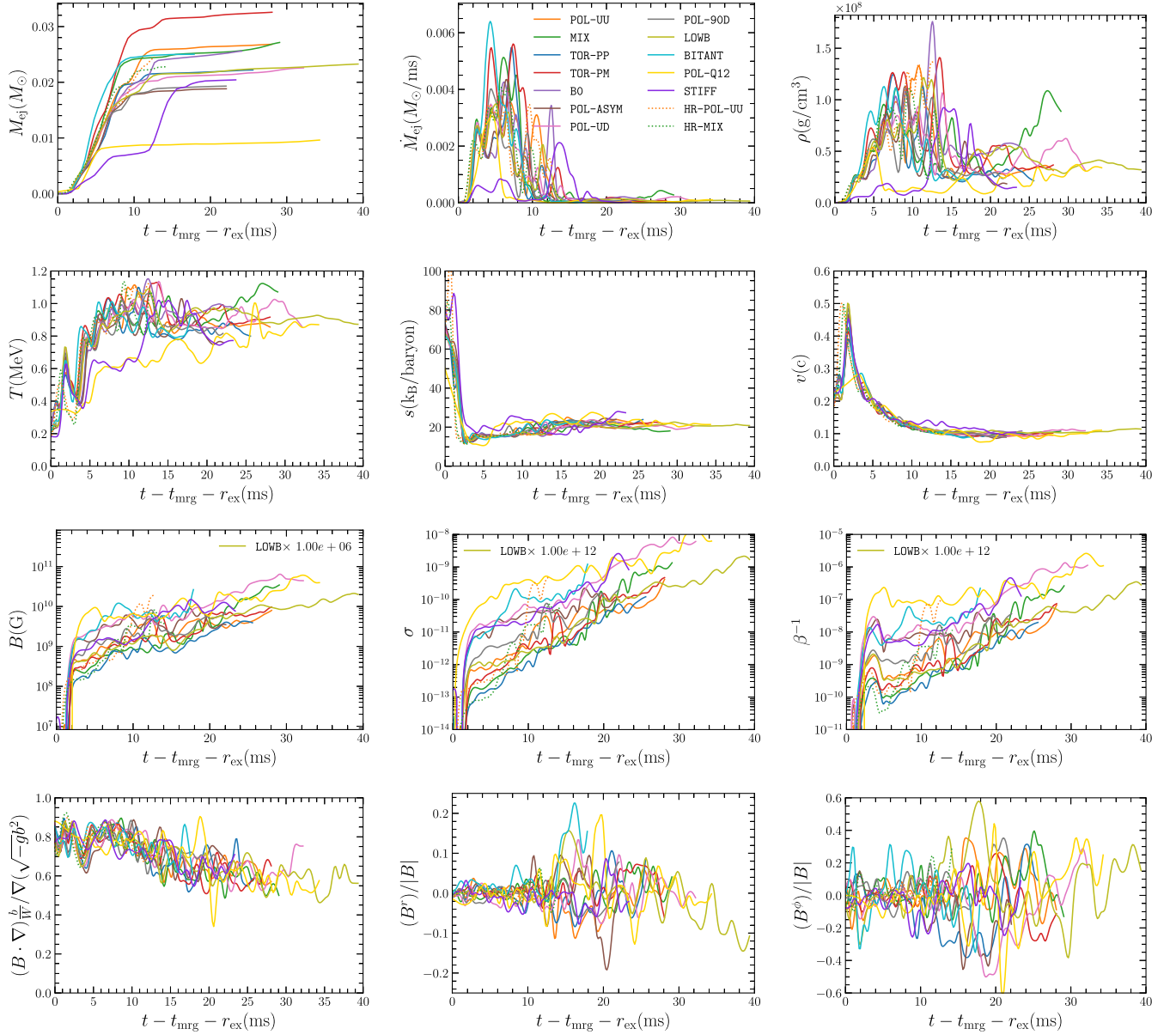


FIG. 10. Upper left panel: the total ejected mass from the binary merger. Upper central panel: the ejected mass flux. Upper right panel: the average ρ of the ejecta. Remaining panels show (from left to right, top to bottom) the mass weighted average of the: temperature, specific entropy per baryon, velocity, magnetic field, magnetization, β^{-1} , ratio between magnetic tension and pressure, ratio between radial magnetic field strength and total magnetic field strength, and ratio between toroidal magnetic field strength and total magnetic field strength. Quantities are plotted as a function of time after merger, less the extraction radius. Note that, when denoted in the legend, the LOWB run has been rescaled for visualization purposes.

We first discuss the total amount of matter ejected from the BNS for the various magnetic configurations studied, demonstrated in the top left panel of Fig 10. All equal mass configurations result in the range $0.019 - 0.033 M_{\odot}$ being ejected. The presence of, and particulars of, the geometry of a magnetic field can have an effect of $\sim 25\%$ on the overall mass ejected, with the unmagnetized case ejecting approximately $0.023 M_{\odot}$. These values are consistent with the values seen by, for instance, [118] for an equal mass configuration, albeit with a different EOS (the “APR4”

model of [119]). More recent studies such as [11,120] see slightly less ejected mass, though the overall hydrodynamical evolution appears consistent, for instance, in the average entropy distribution. We note that these references use an M0 neutrino transport scheme, capturing effects neglected within our simulations.

Considering the purely hydrodynamical nature of the ejecta, we see similar behavior for all magnetic field configurations, with the following exceptions, as shown in Figs. 9 and 10. At late times we see continued emission,

absent for other runs, from MIX and POL-UD, more than 25 ms after merger. These are the configurations that are amplifying the magnetic field fastest at the end of the simulation that have undergone the strongest amplification, respectively, suggesting the impact of the strong magnetic field that develops in the remnant, and associated extra pressure, in driving this ejecta. This is higher density material from the remnant, as shown by the increase in average density in the top right panel of Fig. 10. The velocity and entropy of the ejected material is also largely consistent between runs. We note that there is no clear hierarchy in the amount of ejected matter, the thermodynamics, or the velocity of the ejecta between the B0, LOWB and more strongly magnetized configurations. At longer timescales, with the onset of a magnetically driven wind, we may expect such a hierarchy to form, however. The ejecta is approximately symmetrically distributed over the colatitudinal angle θ with the variation of the mean angle between the configurations 0.1 radians in size, centered on $\frac{\pi}{2}$, with no sign of the imprint of the symmetry breaking in the magnetic field discussed above imprinted onto the ejecta. All configurations show a broadly consistent distribution, with the ejecta peaking in the orbital plane and just above and below, and decaying towards the polar regions, with the origins of these components discussed above.

These results suggest that the difference in magnetic dynamics at the Kelvin-Helmholtz phase are not sufficient to generate pressures strong enough to influence the hydrodynamics of the early ejecta, though the late time fields generated by winding do begin to show signs of increased ejection. The key differences seen in the hydrodynamical properties are seen between different EOS and mass ratios, however. We expect the ejecta from the softer EOS to be dominated on the timescales simulated by shock driven ejecta, due to more compact cores bouncing, and the stiffer EOS to be dominated by tidally driven ejecta [117]. In line with these expected tidal origins of the ejecta, we see that configuration STIFF shows a more extended gradual mass flux, and suppressed temperatures compared to the other configurations, and a more isotropic mass distribution. In line with this expectation we also see less high entropy material with $s > 100k_B/\text{baryon}$ emitted. The unequal mass case POL-Q12 shows the least material ejected overall. Previous results have suggested that for a soft equation of state there is a weak dependence of ejecta mass on the mass ratio, in comparison to the strong enhancement that may be found for a stiffer equation of state [121]. The lack of enhancement in this case for the SFHo EOS is therefore not in conflict with this result. POL-Q12 shows slower ejecta also, with almost no material ejected with velocity over $0.5c$. Dietrich *et al.* [122] have suggested that there should be a weak dependence between the velocity of ejecta and mass ratio, while Vincent *et al.* [123] have found a suppressed velocity for higher mass ratios. Here we find a mild suppression of the

average velocity at higher mass ratio, supporting these observations. The ejecta from the unequal mass case also shows suppressed temperatures, and suppressed entropy, again consistent with being dominated by tidal ejecta.

3. Magnetic field in the ejecta

We now consider the impact of the magnetic field configuration on the magnetic field itself in the ejecta. In the left panel, third row of Fig. 10 we demonstrate the average magnetic field strength. We note that the hierarchy of this strength shares some features with, but does not identically match, that of the global magnification post-merger seen in Fig. 3. Those configurations with the strongest and fastest amplifications such as BITANT, POL-UD, and MIX do see larger corresponding magnetic fields in the ejected material, while the much less magnetized toroidal configurations show lower field strengths. At ~ 25 ms postmerger, we see that the magnetic field in the ejecta may vary by over an order of magnitude, comparing the least strong field TOR-PP and the strongest POL-UD, which, at its strongest, reaches 10^{11} G. Clearly unlike the hydrodynamical evolution, the magnetic evolution of the ejecta is impacted both by the magnetic field dynamics during the KHI driven phase, and the postmerger winding phase, since strong fields are seen in the ejecta for configurations that most strongly amplify in either the former or the latter, e.g., in POL-UD and MIX, respectively.

Other than the STIFF configuration, all models have ejected the majority of the unbound material that they emit over the course of the simulation by ~ 15 ms postmerger (although we may expect on longer timescales more material to be ejected by magnetically driven processes or neutrino driven winds were the effects included here), with the flux dropping considerably after this point. The amplification that we see in the magnetic field, however, continues approximately exponentially after this time, suggesting that the continuous ejecta is being magnetized more strongly by the growing magnetic field in the remnant. This is supported by the central and right plots in the third row of Fig. 10, in which we see a similar hierarchy in the magnetization $\sigma = B^2/\rho$ and β^{-1} . We note that β^{-1} still shows a fluid dominated by fluid pressure, rather than magnetic pressure. The one major difference between these variables and the magnetic field strength can be seen for the STIFF configuration, which has a comparatively amplified σ and β^{-1} at early times, which arises from the comparatively low density and low temperature in the material ejected over the first 10 ms postmerger in this configuration, seen in Fig. 9. The strong magnetization in the ejecta for both EOS suggests that both the tidal and polar shocked ejecta are imprinted with the growing magnetic field of the remnant, as further discussed below.

The impact of the magnetic field on the transport of linear momentum can be separated into two contributions, that of magnetic pressure, proportional to the gradient of

the comoving magnetic field strength $\nabla(\sqrt{-g}b^2)$, which acts to separate tightly spaced magnetic field lines, and that of magnetic tension, proportional to the directional derivative of the magnetic field along the magnetic field lines $(B \cdot \nabla) \frac{b}{w}$, which acts to straighten curved field lines. In the lower left panel of Fig. 10 we show the ratio between these quantities to estimate the more powerful effect in the ejection of unbound material. Here we see that, at early times, the quantities have comparable contributions in the ejecta, but that the contribution of magnetic tension gradually decays postmerger, leaving magnetic pressure the dominant contribution to the ejecting force by a factor of ~ 2 , 30 ms postmerger. This behavior seems largely universal, and independent of the initial magnetic field configuration. Nevertheless, it seems that the tangled nature of the magnetic field means that magnetic tension plays as important role in ejecta evolution, as magnetic pressure. This emphasizes the importance of including fully 3D effects in long-lived simulations of ejecta evolution.

We now discuss the distribution of the magnetic field as a function of the ejecta density, and the angle of emission. As a reference, we demonstrate the full hydrodynamical and magnetic state of the ejected material for the POL-UU configuration as a function of the density and the angle θ in Fig. 11. As the ejecta evolves, the ejected material transitions to larger magnetic fields, associated to the amplification in the remnant, and correspondingly higher magnetizations and values of β^{-1} , in line with the increase in density. At later times we see a steady increase in the matter ejected at higher magnetic field strengths ($\sim 10^{11}$ G) at densities below the maximum, at $\sim 2 \times 10^7$ g/cm³. This leads to the small bump at this density that can be seen in the histograms for these magnetic field quantities, especially σ and β^{-1} . This corresponds to more strongly magnetized matter outflowing from the increasingly magnetized merger remnant, as winding continues, flowing into lower density material. Observing the angular distribution of these magnetic field dependent quantities, we see that the majority of the tidal ejecta, distributed close to the equatorial plane, has approximately constant magnetic field strength, $\sim 10^9$ G, while the increase in magnetic field strength in slightly lower density material observed above occurs in the polar regions, corresponding to the small bump in all magnetic field quantities seen in the polar region above the orbital plane. Comparing with Fig. 7 (upper) we see that shortly postmerger the strongest magnetic fields lie in the equatorial plane, in ejecta driven by rotational effects, while at later times (lower panel), as the toroidal magnetic field builds up and emerges from the remnant (see Fig. 5 of Paper I) the magnetic field begins to emerge in the polar region, emitted from the twin core structures of the rotating remnant at later times, magnetizing the ejecta in this region more strongly.

In Fig. 12 we see the distribution of the magnetic field against the density for all configurations. We see a mild

difference in the lower density tails of premerger unphysical material which we ascribe to differences in reconstruction methods (see Appendix), but overall see the same evolution from low density less magnetized material, ejected early on to higher density more magnetized material in all configurations. In the configurations which show the greatest magnetic field enhancement in the ejecta we see the bump feature previously seen in the POL-UU ejecta clearly enhanced, especially in MIX. This feature appears to be the source of the enhanced magnetic field strength in the MIX configuration, since the majority of the ejecta for MIX is centered around a lower average magnetic field strength than POL-UU, however, this bump reaches higher magnetic field values than seen in POL-UU.

We also note that the high magnetic field strength in the ejecta in the BITANT configuration does not come from the development of such a feature, and simply comes from the overall magnitude of the magnetic field in the ejecta being larger. In Fig. 13 we see the distribution of β^{-1} , which demonstrates a similar hierarchy to that of the magnetic field strength, with a similar bump feature developing below the maximum ρ at high β^{-1} .

In Fig. 14 we demonstrate the angular distribution of the magnetic field, using β^{-1} as a proxy for this (though we note that B shows qualitatively the same behavior). Here we find confirmation that the configurations that have the strongest magnetic fields in the ejecta correspond to those forming a strong field in the polar regions, as can be seen especially in the MIX and POL-UD configurations, with the bump feature corresponding to this strong polar field, supporting the interpretation of Fig. 7.

In Figs. 12 and 13 we have seen the relationship between density and magnetic field strength and β^{-1} . Neglecting the initial low mass unphysical ejecta that arrives at the extraction radius premerger, the ejecta evolution follows the approximately straight line in these log-scaled histograms visible from the bottom left to top right of these plots. Once the initial ejected material has passed through the extraction radius, we see the material slowly begin to move to larger values of the magnetic field as time evolves, as discussed above. To model this evolution we fit a power law relationship of the form

$$B = A_1 \rho^{\alpha_1} \quad (16)$$

$$\beta^{-1} = A_2 \rho^{\alpha_2}, \quad (17)$$

obtaining, at least, squares fit for the parameters A_i and α_i to the mass weighted distribution of ejected material. We perform such fits for all the considered magnetic field configurations, and show the time evolution of the exponent α_i in Fig. 15. Here we see that the magnetic field configuration in the initial star has a nontrivial effect on the magnetic field distribution as a function of the density in the ejecta. Up until ~ 5 ms postmerger our fits are affected

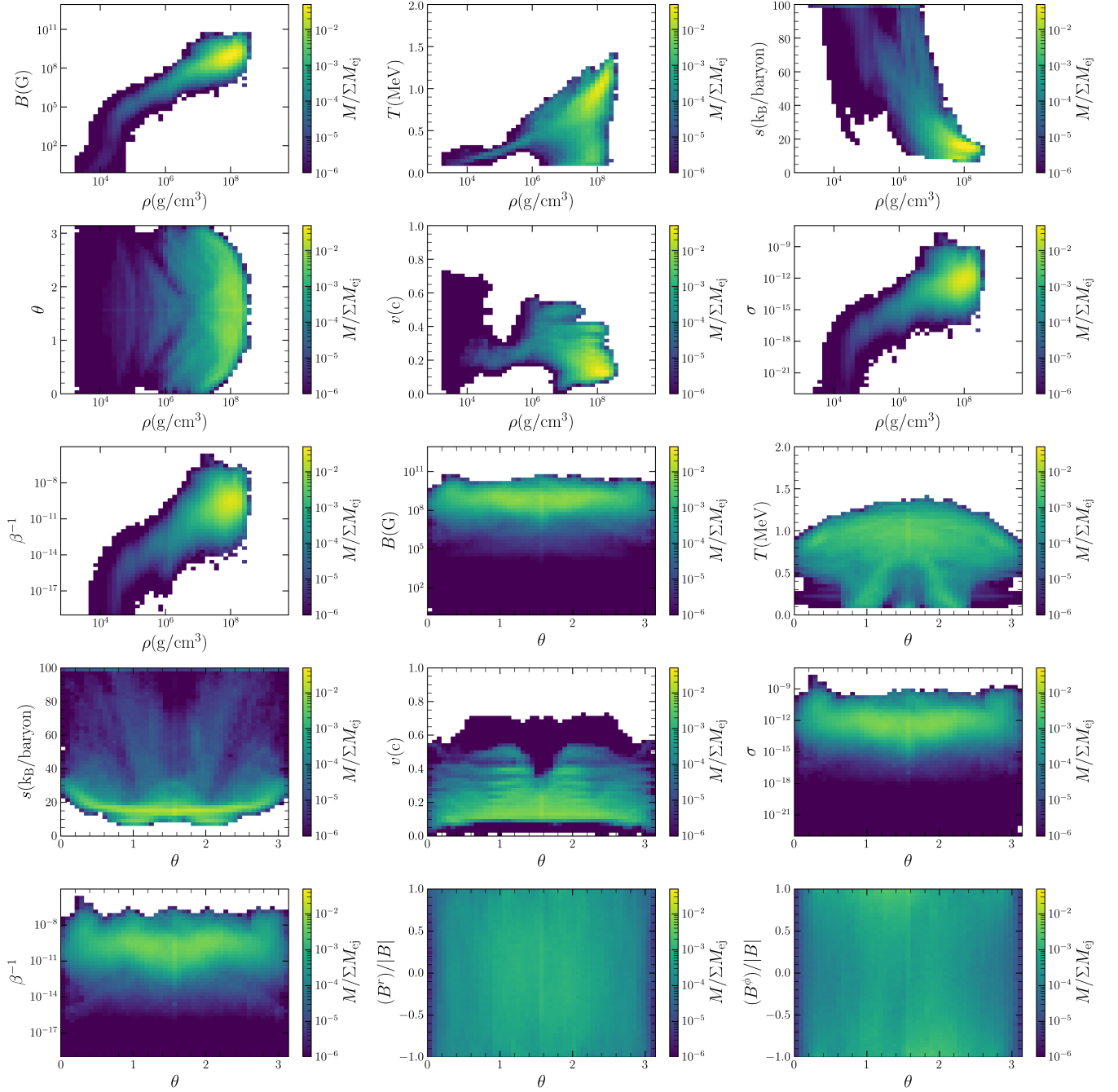


FIG. 11. 2D histograms demonstrating the make up of the cumulative ejected matter at the end of the POL-UU simulation. From left to right, top to bottom, the histograms demonstrate the following pairs of variables on the (x, y) axes: (ρ, B) , (ρ, T) , (ρ, s) , (ρ, θ) , (ρ, v) , (ρ, σ) , (ρ, β^{-1}) , (θ, B) , (θ, T) , (θ, s) , (θ, v) , (θ, σ) , (θ, β^{-1}) , $(\theta, B^r/|B|)$, $(\theta, B^\phi/|B|)$. The color shading demonstrates the cumulative mass ejected normalized by the total mass ejected.

by the unphysical premerger ejecta, but after this point we see that the scaling exponent becomes approximately constant, with a mild decreasing trend, as the magnetic field strength at lower densities gradually increases, as more strongly magnetized material is ejected.

The exponent is inversely correlated with the average magnetic field strength in the ejecta. For the more strongly

magnetized configurations, strong magnetic fields are distributed over the range of high density material that is initially ejected leading to a lower exponent. In contrast, less magnetized configurations, such as the toroidal runs, show less magnetized material, and consequently, a higher exponent. As the magnetic field amplifies and ejection continues over a range of densities, the exponent decreases.

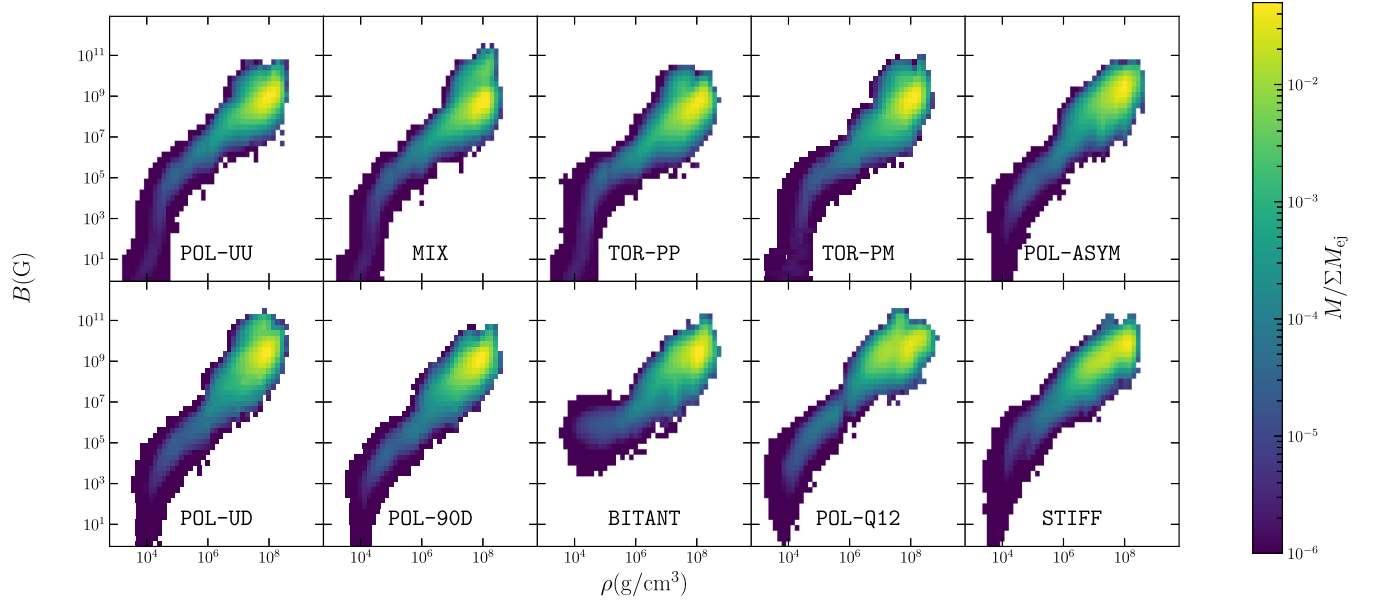


FIG. 12. 2D histograms of magnetic field strength against density for magnetized configurations. The color shading demonstrates the cumulative mass ejected normalized by the total mass ejected for each run.

An overall similar hierarchy is seen for β^{-1} , though with the exponents smaller, demonstrated in the lower panel of Fig. 15. Here the configuration POL-Q12 shows a clearly different behavior to the equal mass configurations, with one of the lowest exponents for the scaling of the magnetic field strength, but the highest for β^{-1} . This feature arises due to the effect of the initial low density ejecta present for the unequal mass configuration, corresponding to the low density excess seen in the density histogram for POL-Q12 in Fig. 9. Equivalently, this can be seen in the

comparatively high value of β^{-1} in the low density tail ($\rho \lesssim 10^6 \text{ g/cm}^3$) of the 2D histogram for POL-Q12 in Fig. 13.

4. Magnetic field structure

We now discuss the magnetic field orientation in the ejected material, by demonstrating the proportion of the magnetic field strength in the radial and toroidal components, respectively, $B^r/|B|$ and $B^\phi/|B|$. The average of the

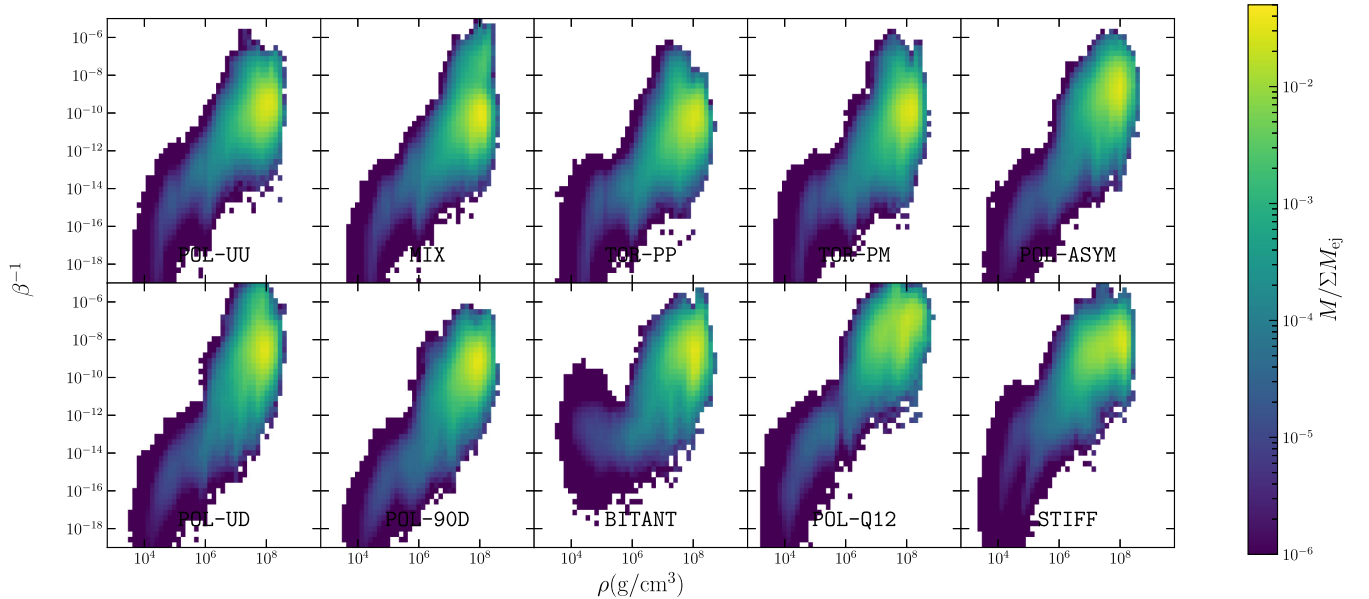


FIG. 13. 2D histograms of β^{-1} against density for magnetized configurations. The color shading demonstrates the cumulative mass ejected normalized by the total mass ejected for each run.

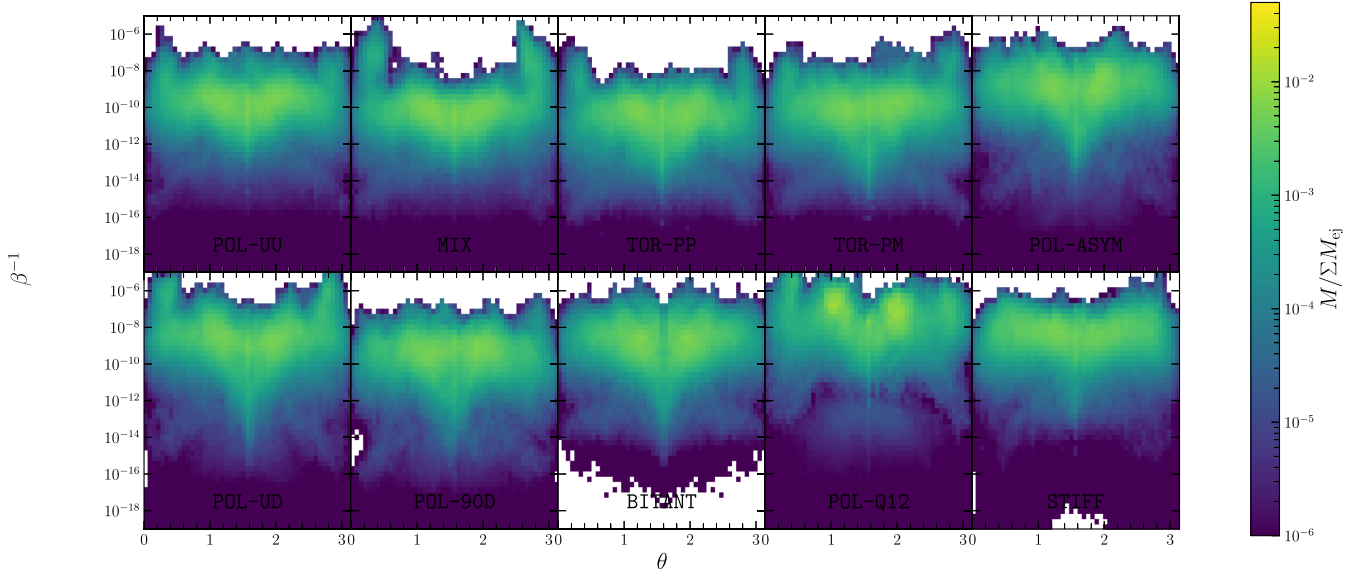


FIG. 14. 2D histograms of β^{-1} against θ for magnetized configurations. The color shading demonstrates the cumulative mass ejected normalized by the total mass ejected for each run.

radial and toroidal components oscillate around zero, with the magnitude of the oscillations, especially for the radial component, strongly increasing after 15 ms postmerger, as seen in the central and right panels of the bottom row of Fig. 10, corresponding to the onset of the winding phase. The size of the oscillations of the toroidal component are notably always larger than that of the radial component, reaching, at their largest, 60% of the total magnetic field strength, while the radial component is mostly below 10%. By the end of the simulation duration considered there is no clear net radial field, nor a global toroidal field present in the ejected material. We see this feature reproduced in the histograms for these quantities, in the right panel of the

third row and left panel of the bottom row of Fig. 9, with the relative strength of the magnetic field in each component shown in the final two panels, with the distributions approximately flat. The majority of the models, with the SFHo EOS and equal mass, show a slightly more toroidally dominated field, with the radial field distribution peaked at 0. The toroidal field contribution is peaked at ± 1 , with the distribution approximately symmetric around 0, other than the POL-Q12 configuration, with mean at -0.09 due to asymmetries in the configuration.

The orientation of the magnetic field in the ejected material may play an important role in its evolution and have a nontrivial impact on the subsequent electromagnetic signature of the kilonova. In, e.g., [72] the magnetic field in the ejected material is modeled with three representative orientations, a purely radial, purely toroidal, and “random” configuration, finding that a toroidal field may “trap” charged particles, giving them more time to thermalize, whereas a radial field may beam these particles outwards, reducing the thermalization of the ejected material. Different choices of such models of the magnetic field may alter the thermalization efficiency of the ejecta significantly and, consequently, impact luminosity predictions for the kilonova light curve. The magnetic field configurations observed here suggest that there is no global radial or toroidal field found in the ejecta and that the random model of [72] is a reasonable choice. To our knowledge this is the first confirmation of this modeling assumption with numerical relativity simulations.

In Fig. 16 we demonstrate the mass averaged relative toroidal component of the magnetic field $B^\phi/|B|$ as a function of θ for the different configurations. Generically across configurations we find a highly oscillatory field as a

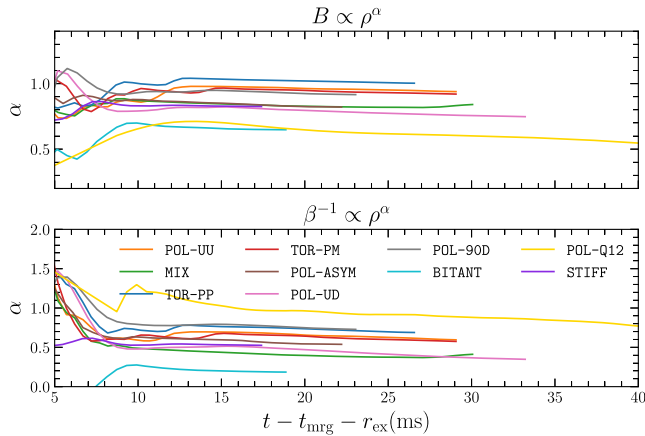


FIG. 15. Upper (Lower) panel: evolution of scaling exponent α_i in ejected material for power law fit $B = A_1 \rho^{\alpha_1}$ ($\beta^{-1} = A_2 \rho^{\alpha_2}$). The evolution of the exponent α is shown as a function of time after merger, less the extraction radius.

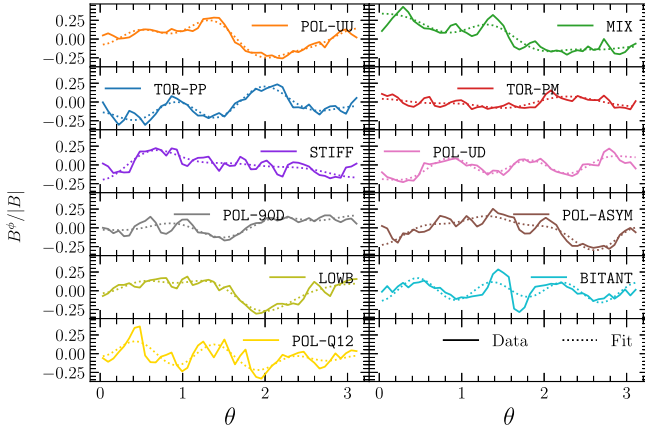


FIG. 16. Mass averaged value of $B^\phi/|B|$ as a function of θ , with truncated Legendre series fit.

function of θ , with the field oscillating between $\sim \pm 40\%$ of the field strength being contained in the toroidal component either co- or counter-aligned with the remnant's rotation. We find the most extreme value for the toroidal field across our simulations in the polar regions, for the MIX configuration, which develops the strongest localized corotating toroidal field component directly above the merger remnant, at $\sim 42\%$ of the total field strength. Here we compare to the results of Paper I, finding that configurations which develop the strongest toroidal components in the remnant appear to also generate the strongest toroidal components in the ejecta in the polar region. This again reinforces the idea of more strongly magnetized remnants building a strong toroidal field which magnetizes ejected material in the polar regions. Configuration POL-UU, as well as those that resemble its magnetic field geometry (MIX, POL-ASYM), shows a sign flip from co- to counter-aligned magnetic fields across the equatorial plane, consistent with the bending of magnetic field lines through winding. Toroidal or antialigned configurations do not show this behavior, suggesting a more tangled geometry in the field lines as the star winds.

Finally we comment on the multipolar structure of this field. We decompose the fraction of the magnetic field in the toroidal component as a Legendre series in $\cos(\theta)$, $B^\phi/|B| = \sum_\ell a_\ell P_\ell(\cos\theta)$. In Table II we list the term in the series corresponding to the maximum coefficient, $\ell_{\max}$, and the associated correlation length of the magnetic field, given by the extraction radius over this value, $L_{\text{corr}} = r_{\text{ex}}/\ell_{\max}$. In Fig. 16 we add the Legendre series fit, keeping terms up to $\ell = 8$, which is larger than $\ell_{\max}$ for almost all configurations. The more symmetric configurations (POL-UU, TOR-PP) have lower $\ell_{\max}$ compared to their asymmetric counterparts (POL-UD, TOR-PM, POL-Q12), suggesting asymmetry between the stars generates a less coherent field, though we note that the enforced symmetry over the z axis for the BITANT configuration leads to a higher $\ell_{\max}$ than in the POL-UU case.

E. Summary and discussion

We now briefly summarize our results and discuss the context of this work. Our simulations have shown that the presence of magnetic fields enhances the excitement of higher order modes in the postmerger GW spectrum, and that the evolution of these modes depends on the initial magnetic field topology. This excitement is driven by the different distribution of energy density and magnetic pressures at merger. Recent studies have investigated the impact of magnetic field topology on the gravitational wave spectrum [62,63], with simulations at higher resolution than those considered here. These studies have suggested that magnetic fields may have an impact on the postmerger GW signal that is as large as effects from the EOS, which may be detectable in next generation GW detectors. Our error analysis suggests that, for our simulations, at a fixed resolution, the systematic error associated to the hydrodynamical scheme employed is larger than a detectable signature. This suggests that extensive numerical experiments and higher resolution simulations may still be necessary to investigate the detectability of such features. This supports also the claims of Zappa *et al.* [124] who find that multimessenger observables extracted from the postmerger phase of such simulations are incredibly sensitive to the numerics and physics input of a given numerical simulation, as well as its resolution.

We have seen that antialigned magnetic fields in the initial stars strongly enhance the magnetic field amplification from the Kelvin-Helmholtz instability, in line with predictions from local simulations [111], and that toroidal fields can suppress the enhancement, in agreement with other BNS simulations exploring similar parameter choices [60,62]. In long-lived simulations, it has been suggested [68,125] that at late times postmerger the magnetic field should reach a quasiuniversal end state, independent of the initial magnetic field geometry. We have seen in our simulations that the initial magnetic field geometry can have a nontrivial impact on the magnetic field evolution in the ~ 30 ms postmerger, which does not suggest a universal behavior. Longer simulations, however, will be required to test the universal nature of the field at the later times considered in these works, as well as the potential inclusion of a subgrid model to effectively capture the impact of the full cascade of the magnetic field to small scales. This work, however, may support results from simulations performed by Chabanov *et al.* [67] and Jiang *et al.* [126] that suggest that the specifics of the magnetic field strength and geometry may have lasting imprints in a remnant object, in particular, the depth of the magnetic field in the star, as shown by the weaker amplification of the toroidal field configurations.

We have analyzed the evolution of asymmetry in the magnetic field in our simulations and have found that, in the absence of an enforced symmetry, asymmetry will grow at a universal rate regardless of the initial condition chosen.

This suggests the presence of a symmetry breaking bifurcation in the solution that is suppressed when a bitant symmetry is enforced, suppressing turbulent features near the symmetry plane, and overestimating the magnetic field amplification through winding as discussed in Paper I. BNS simulations frequently employ this symmetry to reduce the cost of the simulation, but we have shown here that, in the presence of magnetic fields, this can lead to a significant error in the simulation, with asymmetries growing as large as the symmetric part of the magnetic field.

The different magnetic field configurations considered have shown minimal impact on the hydrodynamical evolution of the ejecta from the BNS, but those configurations that show the strongest global magnetic amplification appear to show early signs of a later magnetically driven ejecta focused in the polar regions. This aligns with longer term simulations performed by e.g., [13], where strong magnetic fields in the postmerger phase begin to drive late time magnetic winds. In [9,12] it was found that strong magnetic fields in the postmerger remnant can impact the hydrodynamical evolution of the ejecta over the timescales we consider here, with a clear dependence on the strength of the magnetic field, which we have not seen. In these simulations, however, a strong poloidal magnetic field is imposed on the postmerger remnant. We have not seen the emergence of a global coherent poloidal field by the end of our simulations, and so we may conclude that the impact of the magnetic field on the hydrodynamics of the ejecta may only be seen after sufficient time has passed to allow this field to develop, or at sufficiently high resolutions capable of resolving the relevant dynamo processes which may drive its formation. As in the remnant, magnetic field topology does have an impact on magnetic field strength in the ejecta, with more magnetized remnants generating more strongly magnetized ejecta. As these more magnetized runs evolve, we see highly magnetized material emitted in the polar regions, driven by the strong toroidal field in the remnant, continually amplifying by winding.

Finally we have seen that magnetic fields are largely randomly oriented in the ejecta, agreeing with modeling assumptions used in e.g., [72]. This suggests that charged particles in realistic ejecta may be trapped closer to the remnant by the magnetic field, giving them more opportunity to thermalize with their environment. A stronger radial field would instead beam these particles outwards, giving them less time to thermalize. This validates estimates of the thermalization timescales of ejecta used in the above work, suggesting that the kilonova modeling associated to this analysis is robust.

IV. CONCLUSION

In this work we have performed a series of simulations of magnetized BNS mergers with varying initial choices of magnetic field topology, also varying the EOS, mass ratio and symmetry assumptions of our numerical domain. The

impact of the choice of initial magnetic field topology has been found to have a significant effect on the magnetic field evolution during the merger, as well as in the ejected matter from the BNS merger.

Firstly we have analyzed the impact of different magnetic field configurations on the GW signal of the BNS. The presence of a magnetic field has minimal impact on the wave signal during the inspiral, but in the postmerger phase, different magnetic field configurations can have an impact on the evolution of GW modes, compared with each other and with a reference unmagnetized configuration. We find that the dephasing between the dominant (2, 2) mode and leading subdominant (2, 1) and (3, 3) modes may be affected by different magnetic field configurations. This may lead to a change in the size of the dephasing that is tens of radians in magnitude, but we note that this effect is highly sensitive to the reconstruction routine used and that the overall magnitude of this effect seen in the (3, 3) mode is less than the difference in dephasing that arises from changing reconstruction routine.

Subsequently we have investigated the impact of the initial field configuration on the magnetic field amplification. Here we have found a large impact, finding that the magnification arising during the KHI may be strongly amplified by the presence of antialigned magnetic fields, with antialigned poloidal and toroidal fields magnifying considerably more than their aligned counterparts, in agreement with local simulations. This feature may arise due to an increase in magnetic tension in the antialigned case. The amplification during this phase for toroidal fields, however, is strongly suppressed, due to the weaker magnetic field strengths in the region where the KHI is present. This suggests that features associated to strong magnetic field development in the postmerger phase, such as the launching of a jet, may be delayed for models of neutron stars with strong toroidal fields, but that magnetic field amplification may be underestimated for generic configurations with different alignments of poloidal fields.

After the KHI has developed, we have also found that the magnetic field topology has a large impact on the amplification due to winding, with the amplification slower in toroidally dominated configurations, but enhanced in the presence of asymmetries, and faster at late times for a combination of poloidal and toroidal fields that resemble the end state of the instabilities that a purely poloidal field will undergo. Enforcing a bitant symmetry also enhances this amplification. We find that, as a consequence, the Poynting flux from the merger remnant is enhanced for those configurations with larger and faster amplifications during the winding phase.

We have analyzed the asymmetries that develop in our simulations when evolving symmetric initial data. We find that, independent of the specific topology of the initial data, the size of the asymmetric magnetic field evolves at an approximately universal rate, saturating at approximately

10% of the symmetric magnetic field in size. This suggests the presence of a spontaneous symmetry breaking bifurcation in the solution to the GRMHD equations, a feature that will be neglected when enforcing a bitant symmetry in such simulations.

We perform a detailed analysis of the ejected material from these simulations. For identical initial hydrodynamical configurations we find that the evolution of the hydrodynamics of the ejected material is broadly similar for different magnetic configurations. The nature of the magnetic fields within the ejecta, however, can be strongly affected by the magnetic field configuration and amplification. We find that the magnetic field strength in the ejecta may vary by over an order of magnitude depending on the magnetic field initial data, with consequent impacts on the β^{-1} and magnetization of the ejected material. We find the strongest magnetic fields in the ejecta at late times in the most strongly amplified configuration POL-UD, as well as the fastest growing MIX configuration. These configurations most readily eject more highly magnetized material in the polar regions with comparatively larger magnetic fields up to 10^{11} G, as a strong toroidal field develops in the remnant. We model the relationship between density and magnetic field strength, and density and β^{-1} in the ejected material, and find a simple power law scaling that decays gradually over time as more, less dense, highly magnetized material begins to be ejected in the polar regions. An enforced bitant symmetry forces a much weaker relationship between the density and magnetic field strength, than the equivalent full grid configuration, and leads to an enhanced field strength also in the ejecta.

We have also analyzed the magnetic field orientation in the ejected material in order to validate existing models used to estimate thermalization efficiencies of charged particles in the ejecta. We find that overall no dominant global field has formed in the ejecta by the end of our simulations and that, consequently, an approximately random model may provide a reasonable approximation to the field. This model increases the timescales over which charged particles may remain trapped in the magnetic field and given an opportunity to thermalize, impacting expected light curve luminosities in the kilonova, compared to a radial field. Towards the end of the timescales considered, however, we begin to see signs of a stronger toroidal field forming, especially for the MIX configuration, in the polar region.

A fuller analysis of the properties of such configurations will require additional features not present in these simulations. Firstly the impact of neutrino effects must be included through, for instance, an M1 scheme [127–129] in order to obtain a more accurate model of the ejecta temperature, entropy and, key for nucleosynthesis estimates, the composition. Longer term simulations will be required to investigate the full impact of the magnetic field configuration on jet launching, and the use of subgrid modeling [130–132] may be key to accurately capturing the

magnetic field in the remnant, and hence the later time magnetically driven ejecta. The impact of the magnetic field on the postmerger GW signal has been shown to be highly dependent on the choice of reconstruction method. Future simulations of these effects should probe even higher resolutions, to investigate the robustness of this feature, and its potential detectability.

ACKNOWLEDGMENTS

E. G. thanks Luciano Combi for helpful discussions. E. G. and D. R. acknowledge funding from the National Science Foundation under Grants No. AST-2108467 and No. PHY-2407681. E. G. acknowledges funding from an Institute for Gravitation and Cosmology fellowship. D. R. acknowledges support from the Sloan foundation and from the National Science Foundation under Grant No. PHY-2020275. J. F. and D. R. acknowledge U.S. Department of Energy, Office of Science, Division of Nuclear Physics under Awards No. DE-SC0021177 and No. DE-SC0024388. S. B., B. D. and M. J. acknowledge support by the EU Horizon under ERC Consolidator Grant No. InspiReM-101043372. S. B. and M. J. acknowledge support from the Deutsche Forschungsgemeinschaft (DFG) under MEMI (BE 6301/2-2, Project No. 443239082). P. H. acknowledges funding from the National Science Foundation under Grant No. PHY-2116686. The numerical simulations were performed on TACC's Frontera (NSF LRAC allocation PHY23001) and on Perlmutter. This research used resources of the National Energy Research Scientific Computing Center, a DOE Office of Science User Facility supported by the Office of Science of the U.S. Department of Energy under Contract No. DE-AC02-05CH11231. Simulations were also performed on the national HPE Apollo Hawk at the High Performance Computing Center Stuttgart (HLRS). The authors acknowledge HLRS for funding this project by providing access to the supercomputer HPE Apollo Hawk under the Grants No. INTRHYGUE/44215 and No. MAGNETIST/44288. Simulations were also performed on SuperMUC_NG at the Leibniz-Rechenzentrum (LRZ) Munich. The authors acknowledge the Gauss Centre for Supercomputing e.V. [133] for funding this project by providing computing time on the GCS Supercomputer SuperMUC-NG at LRZ (allocations pn67xo, pn76li, pn68wi and pn36jo). Postprocessing and development run were performed on the ARA cluster at Friedrich Schiller University Jena. The ARA cluster is funded in part by DFG Grants No. INST 275/334-1 FUGG and No. INST 275/363-1 FUGG, and ERC Starting Grant, Grant agreement No. BinGraSp-714626.

DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not

technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

APPENDIX: EFFECT OF RECONSTRUCTION SCHEME

We have performed simulations with two different reconstruction schemes, PPM and WENOZ. To benchmark the expected numerical errors introduced by varying the reconstruction scheme, we perform the same run, POL-UU, with both schemes. Here we demonstrate the GW signal, magnetic field amplification and ejecta magnetization for both runs to quantify the variation due to reconstruction scheme.

We first note that the merger time is, as expected, delayed for the WENOZ run compared with the PPM run due to a reduced dissipation, with the merger occurring ~ 0.9 ms later in the WENOZ run. Accounting for the shifted time of merger, we see in Fig. 17 the evolution of the (2, 2) mode of the GW strain is almost identical up to merger in both amplitude and instantaneous frequency. In the postmerger phase, we see small deviations in the (2, 2) mode amplitude, and slightly larger oscillations in the frequency for the WENOZ configuration. In the subdominant modes, however, here focussing on the (2, 1) and (3, 3) modes, we see a significant difference in the postmerger evolution. Immediately postmerger we see a large peak in the amplitude of the (2, 1) mode in the WENOZ configuration, a factor 2 larger than that of the PPM configuration, with later oscillations in the PPM configuration larger than those in the WENOZ configuration. Conversely in the (3, 3) mode, we see two peaks in amplitude immediately postmerger for the PPM configuration. These peaks are approximately coincident with, but larger than, corresponding peaks in

the WENOZ configuration, with the initial peak a factor 4 larger in the PPM configuration. After these peaks we see the WENOZ amplitude growing above the PPM values. For the instantaneous frequencies of the (2, 1) and (3, 3), the inspiral is dominated by noise, due to the overall small value of the amplitude of these modes, but in the postmerger, other than isolated spikes in the frequency, the overall frequency evolution is broadly similar between the reconstruction methods. To estimate the error associated to the choice of reconstruction scheme on the GW dephasing discussed in Sec. III A, we calculate the accrued dephasing $\Delta\phi_{21}$ and $\Delta\phi_{33}$ for the two POL-UU configurations. For the PPM (WENOZ) configuration, the accrued dephasing is 29 (17) radians for the (2, 1) mode and -30 (-48) radians for the (3, 3) mode. This gives an error of 12 and 18 radians, respectively. These are large errors compared to the maximal size of the dephasing effects seen in Sec. III A, of size 39% and 123%, respectively, though we note that it is still the case that for a consistent choice of reconstruction scheme, a range of dephasings of size up to 12 and 8 radians can be found for the (2, 1) and (3, 3) modes, respectively, for different choices of magnetic field data. The size of these errors emphasizes the importance of robustly exploring the impact of not just resolution but also the choice of numerical scheme used to evolve the hydrodynamics, on any potentially detectable signal found in the postmerger GW.

We now investigate the impact of reconstruction on the magnetic field amplification, demonstrated in Fig. 18. Again accounting for the different merger time, we find that the magnetic field amplification is relatively robust to the choice of reconstruction method. When drawing a comparison with the reconstruction analysis of [50], we see that the WENOZ method should preserve the central density of the star more accurately than the PPM method, but the star surface may be preserved less well. This may lead to a slightly more diffuse outer layer of the star, where we expect the KHI to develop. Therefore the development of this feature may be dependent on the impact of the

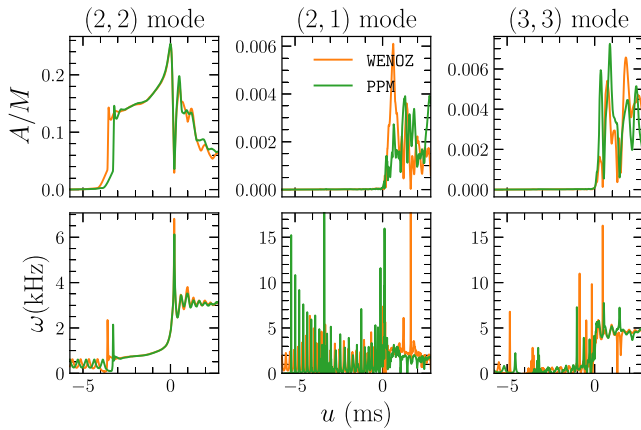


FIG. 17. Upper row: the GW strain amplitude for configuration POL-UU with different reconstruction routines. Bottom row: the instantaneous GW frequency. Left, central and right columns show the $(\ell = 2, m = 2)$, $(2, 1)$, $(3, 3)$ modes, respectively.

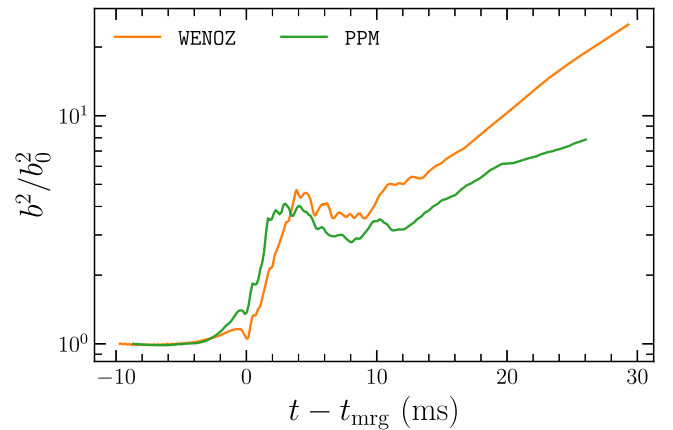


FIG. 18. The magnetic field amplification for configuration POL-UU with different reconstruction routines.

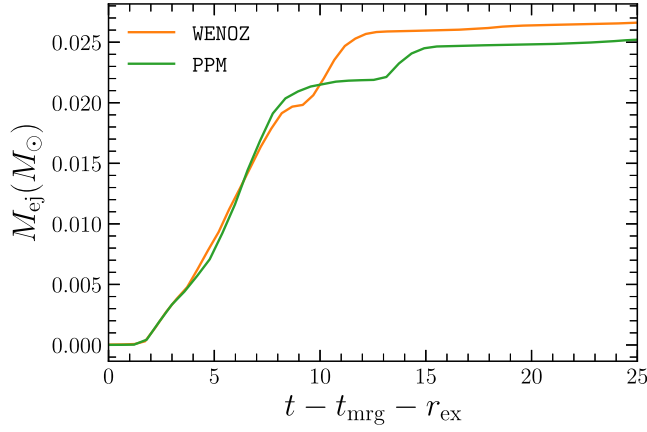


FIG. 19. The ejected mass for configuration POL-UU with different reconstruction routines.

reconstruction scheme on the hydrodynamical evolution. We find that the PPM configuration shows an earlier start to the magnetic field amplification, and that the saturation of the KHI occurs more quickly than the corresponding WENOZ configuration, which we attribute to the merger process happening slightly quicker. Once the KHI has saturated, the amplification of the magnetic field between the two runs is very similar, with the WENOZ run amplifying by a factor of $\sim 10\%$ more than the PPM run. After the saturation of the KHI, the intermediate phase progresses similarly for both configurations, with the magnetic field strength remaining constant, and lasting for approximately the same duration. As the winding phase begins, the slightly larger magnetic field for the WENOZ configuration results in a larger field generated by the end of the runs by a factor of ~ 2 . We note that, while this is again a large factor if one wishes to make precise quantitative predictions from such simulations, a factor of 2 is subdominant to the 2 orders of magnitude seen across different magnetic field initial data configurations in Fig. 3, and so we consider this analysis robust to choice of reconstruction method.

Finally, we comment on the impact of reconstruction scheme on the ejecta, by comparing the total ejected mass, and the average magnetic field strength in the ejecta. In Fig. 19, for the PPM run we see that initially after merger the first initial flux of ejecta is larger than the corresponding WENOZ run, leading to a larger mass ejection, but that subsequent ejections, corresponding to core bounces, are

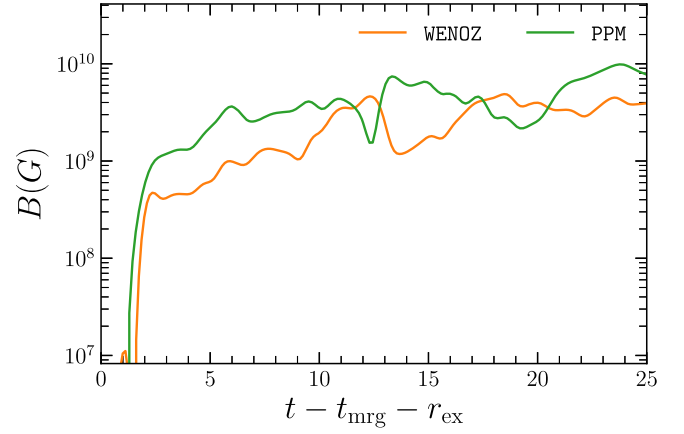


FIG. 20. The magnetic field strength in the ejected material for configuration POL-UU with different reconstruction routines.

less extreme than in the WENOZ configuration, and consequently, marginally more mass is ejected for the WENOZ run. At the end of the simulations the WENOZ run has ejected $\sim 0.027 M_\odot$, while the PPM configuration ejects $\sim 0.025 M_\odot$, a 7% error.

The magnetic field strength in the ejecta can also be affected by reconstruction scheme choice. In Fig. 20 we find that, contrary to the hierarchy during the winding phase, the PPM configuration provides slightly more magnetized ejecta, at least at early times postmerger. At ~ 5 ms postmerger we see that the PPM configuration has on average a magnetic field ~ 4 times stronger than in the WENOZ configuration. By the end of the simulations, however, we see that the two configurations have magnetic fields of approximately equal magnitude on average, oscillating around a common value. The main impact on the ejecta from the reconstruction method appears to be on the unphysical low density material arriving at the extraction spheres before merger time. For the WENOZ configuration we see more low density material arriving at early times, which appears to be absent for the PPM configurations, which, again, supports the analysis of reconstruction methods in [50]. These features can be seen in the low density tails of the 1D and 2D histograms in Sec. III D, but we note that these tails contain very little mass, and that the bulk of the physical ejecta is comparatively unaffected, as justified by the comparatively small errors found in this direct comparison.

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