

Gamma-ray bursts and kilonovae from the accretion-induced collapse of white dwarfs

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ABSTRACT

We present the first seconds-long 2D general relativistic neutrino magnetohydrodynamic simulations of accretion-induced collapse (AIC) in rapidly rotating, strongly magnetized white dwarfs (WDs), which might originate as remnants of double-WD mergers. This study examines extreme combinations of magnetic fields and rotation rates, motivated both by the need to address the limitations of 2D axisymmetric simulations and to explore the physics of AIC under rare conditions that, while yet to be observationally confirmed, may be consistent with current theoretical models and account for unusual events. Under these assumptions, our results demonstrate that, if realizable, such systems can generate relativistic jets and neutron-rich outflows with properties consistent with long gamma-ray bursts (LGRBs) accompanied by kilonovae, such as GRB 211211A and GRB 230307A. These findings highlight the potential role of AIC in heavy *r*-process element production and offer a framework for understanding rare LGRBs associated with kilonova emission. Longer-duration 3D simulations are needed to fully capture magnetic field amplification, resolve instabilities, and determine the fate of the energy retained by the magnetar at the end of the simulations.

1. INTRODUCTION

The nature of the central engines driving long-duration gamma-ray bursts (LGRBs) accompanied by kilonova emissions, such as GRB 211211A (Rastinejad et al. 2022; Troja et al. 2022; Mei et al. 2022; Yang et al. 2022) and GRB 230307A (Fermi GBM Team 2023; Dalessi & Fermi GBM Team 2023; Dalessi et al. 2023; Evans & Swift Team 2023; Burrows et al. 2023), remains unresolved. These events have been suggested to originate from neutron star mergers (Gal-Yam et al. 2006; Rastinejad et al. 2022; Troja et al. 2022; Yang et al. 2022). While neutron star mergers are known to power kilonovae via the rapid neutron capture (*r*-process) nucleosynthesis (Li & Paczyński 1998; Metzger et al. 2010; Tanaka 2016; Metzger 2019), the long-duration GRB is

difficult to reconcile with ($\lesssim 1$ s) cooling timescale of the remnant disk formed in these systems (e.g., Fernández & Metzger 2016). See however Gottlieb et al. (2023) for a compact-merger-powered long-GRBs model. An alternative scenario to explain these events includes the collapse or merger of white dwarfs (WDs; Fryer et al. 1999b; Metzger et al. 2008; Chen et al. 2024). However, such models have not been studied numerically as extensively as compact binary mergers.

WDs are compact objects supported by electron degeneracy pressure. The fate of an accreting WD depends on its mass, composition, and accretion rate (Nomoto 1986). Low-mass carbon-oxygen WDs ($\lesssim 1.2 M_{\odot}$) with high accretion rates are likely to undergo carbon deflagration and explode as Type Ia supernovae, leaving no remnant (Yoon et al. 2007; Shen & Bildsten 2009; Shen et al. 2009; Moore et al. 2013). In contrast, more massive oxygen-neon (O+Ne) WDs ($\gtrsim 1.2 M_{\odot}$) with slower accretion rates can result in higher mass WDs. When the total gravitational mass of a WD exceeds the Chandrasekhar limit of approximately $1.44 M_{\odot}$ (Chandrasekhar 1931), gravitational collapse into a neutron

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star could take place, resulting in AIC event (Nomoto & Kondo 1991).

The progenitors of AICs are expected to rotate rapidly due to their accretion history (Piersanti et al. 2003a,b; Uenishi et al. 2003; Saio & Nomoto 2004; Kashyap et al. 2018), and are likely strongly magnetized (Ferrario et al. 2015; Kawka 2020; Ferrario et al. 2020; Zhu et al. 2015; Pakmor et al. 2024a). Such collapses are expected to result in dim and fast evolving transients, and have been suggested as possible sources for next-generation ground-based gravitational-wave observatories (Dimmelmeier et al. 2008; Longo Micchi et al. 2023). Moreover, they might power gamma-ray bursts (Yi & Blackman 1998; Metzger et al. 2008; Perley et al. 2009) and fast radio bursts (Waxman 2017; Margalit et al. 2019). These collapses may further result in the formation of a millisecond pulsar or low-mass X-ray binary (Ablimit & Li 2015; Wang et al. 2022; Ablimit 2023), and could also be sites of production of rare neutron-rich heavy elements via the r -process nucleosynthesis (Fryer et al. 1999a).

Despite the significant astrophysical potential of AICs, the current understanding primarily stems from either non-magnetized or Newtonian simulations (Baron et al. 1987; Woosley & Baron 1992; Fryer et al. 1999a; Dessart et al. 2006, 2007; Abdikamalov et al. 2010; Sharon & Kushnir 2020; Longo Micchi et al. 2023; Mori et al. 2023). To date, no studies have employed general relativistic neutrino magnetohydrodynamics (GR ν MHD) to explore AICs of WDs. In this work, for the first time, we perform axisymmetric GR ν MHD simulations of rapidly rotating, magnetized WDs with energy-integrated neutrino transport, focusing on jet formation and matter outflows. The paper is organised as follows. In section 2, we outline the methods we used in this work. We present our results in section 3, and summarise our findings and conclusions in section 4.

2. METHODS

The primary goal of our study is to explore whether extreme cases of AIC can drive relativistic jets and produce GRBs, which are much rarer than Type Ia SNe. Hence, we focus on events where extreme initial conditions may be critical. Our simulations model a specific subclass of progenitors: highly magnetized, rapidly rotating super-Chandrasekhar WDs formed through the AIC process. These WDs differ fundamentally from typical, stable, and observationally constrained WDs. The proposed progenitors arise under extreme conditions, which may naturally result from various AIC formation channels, including single-degenerate or double-degenerate WD mergers. These scenarios are conducive to substantial magnetic field amplification during accretion or merger phases, potentially driven by dynamo mechanisms linked to the magneto-rotational instability (e.g., Zhu et al. (2015); Pakmor et al. (2024b)). Moreover, the rapid rotation assumed in our mod-

els aligns with the expected properties of these super-Chandrasekhar WDs, which may rotate at or beyond the Keplerian limit. Consequently, our simulations focus on rapidly rotating, strongly magnetized super-Chandrasekhar WDs as potential progenitors for AICs leading to GRBs.

The initial equilibrium configuration of the rapidly rotating relativistic progenitor used in our simulations are generated with RNS code (Stergioulas & Friedman 1995). We adopt the “LS220” equation of state (EOS; Lattimer & Swesty 1991). The initial configuration is constructed at a fixed temperature of 0.01 MeV and an electron fraction $Y_e = 0.5$. The WD has an initial mass of $1.5 M_\odot$, with a central energy density $\epsilon_c/c^2 = 1 \times 10^{10} \text{ g} \cdot \text{cm}^{-3}$ (Nomoto & Kondo 1991; Yoon & Langer 2004, 2005; Dessart et al. 2006, 2007). The degree of stellar rotation is characterized by the ratio between the polar and equatorial radii, a_r , which reflects the star’s rotational deformation. Since progenitors of AICs are expected to rotate rapidly due to their accretion history (Piersanti et al. 2003a,b; Uenishi et al. 2003; Saio & Nomoto 2004; Kashyap et al. 2018), we focus on rapidly rotating WDs with $a_r = 0.75$, corresponding to an angular velocity $\Omega \approx 5 \text{ Hz}$, which is $\sim 80\%$ of the Kepler limit.

We endow the initial WD profile with a purely poloidal magnetic field that is roughly uniform inside a sphere of radius $r_0 = 600 \text{ km}$ (Dessart et al. 2007), where the poloidal component is aligned with the axis of rotation. Such magnetic field is described by the following vector potential (e.g., Suwa et al. 2007; Varma et al. 2021):

$$\left(A^{\hat{r}}, A^{\hat{\theta}}, A^{\hat{\phi}} \right) = \frac{r_0^3}{2(r^3 + r_0^3)} (0, 0, B_{\text{pol}} r \sin \theta), \quad (1)$$

where r is the distance from the center. To account for the instabilities that can amplify the magnetic field, and that cannot be resolved due to axisymmetry and limited resolution of our simulations, we consider large initial magnetic fields up to 10^{12} G . Specifically, we perform simulations with a purely poloidal magnetic field, using four initial field strengths: $B_{\text{pol}} = \{10^9, 10^{10}, 10^{11}, 10^{12}\} \text{ G}$. Also, we superimpose a temperature profile $T = T_c (\rho/\rho_c)^{0.35}$, with $T_c = 5 \times 10^9 \text{ K}$ ($\approx 0.43 \text{ MeV}$) similar to Dessart et al. (2006, 2007) before the dynamical simulations.

Our models are evolved with the GR ν MHD code Gmumu (Cheong et al. 2020, 2021, 2022, 2023; Ng et al. 2024), which solves GR ν MHD and Einstein field equations in the conformally flat approximation, and uses the energy-integrated two-moment radiation transport scheme of Radice et al. (2022) for neutrino transport. Neutrino rates are provided by coupling WeakRates module in WhiskyTHC code (Radice et al. 2022). The divergence-free condition of the magnetic field is preserved by using staggered-meshed constrained transport (Evans & Hawley 1988).

All simulations are axisymmetric and performed in cylindrical coordinates (R, z) , with a computational domain extending from $0 \leq R \leq 2000$ km and $0 \leq z \leq 2000$ km. The base resolution is $n_R \times n_z = 128 \times 128$ grid points in R and z , with six levels of adaptive mesh refinement applied. At the center of the star, the finest grid resolution is $\Delta R = \Delta z \approx 488$ m when the maximum density is beyond 10^{12} g · cm $^{-3}$.

Our simulations adopt the Harten, Lax and van Leer (HLL) approximated Riemann solver (Harten et al. 1983), the 3rd-order reconstruction piecewise parabolic method (PPM) (Colella & Woodward 1984) and the IMEXCB3a time integrator (Cavaglieri & Bewley 2015). The finite temperature equation-of-state “LS220” (Lattimer & Swesty 1991) is used for the evolution. To capture the dynamics of the low density gas, we extend the equation-of-state by assuming the gas is a mixture of electrons, ions, and photons. The details of this extension are presented in the appendix A.

Energy-coupled neutrino interactions such as neutrino-electron inelastic scattering, which are important in the collapsing phase, are not included when an energy-integrated neutrino transport is used. An effective approach to overcome this issue is to apply the parametrised deleptonisation scheme of Liebendörfer (2005) for electron fraction evolutions. We switch to the two-moment neutrino transport when the core bounces, and enable the coupling between neutrinos and the fluid, as described in Cheong et al. (2024).

We investigate the kilonova emission starting from the ejecta profile extracted from the simulation with the strongest magnetic field, $B_{\text{pol}} = 10^{12}$ G. Unbound material is identified via the Bernoulli criterion and collected at an extraction radius of $r_{\text{ext}} = 1800$ km. To account for the angular dependence of the ejecta properties and model the long-term evolution, the polar dependence is discretized into 90 angular sections. Thermodynamic properties such as entropy, temperature, and electron fraction for each angular section are tabulated as a function of the enclosed ejecta mass using a mass-weighted average, procedure adopted also in (Wu et al. 2022; Magistrelli et al. 2024). The time-dependent ejecta profile is mapped into a Lagrangian profile, with the latest shell at the extraction radius and previous layers placed on top. We solve the spherically symmetric hydrodynamic equations using a ray-by-ray approach, treating each angular section as a 1D problem. The mass of each shell is scaled by the factor $4\pi/\Delta\Omega$, where $\Delta\Omega = 2\pi \sin\theta d\theta$ represents the solid angle of the section. The position $r(m)$ of each mass shell is computed based on the constraint $m(r) = 4\pi \int_{r_{\text{ext}}}^r \rho(r) r^2 dr$, where ρ is the mass density and $m(r)$ is the enclosed mass. We neglect non-radial flows of matter and radiation between different sections. The ejecta profile is evolved using the 1D radiation hydrodynamics code SNEC as implemented in Morozova et al. (2015) and Wu et al. (2022). Kilonova light curves are calculated using the analytic, time-

independent opacity model from Wu et al. (2022). The results are mapped back into an axisymmetric framework, rescaling the luminosities by $\Delta\Omega/4\pi$, and combining them to calculate the final kilonova emission.

3. RESULTS

Collapse dynamics—The WD collapses due to electron capture and the resulting deleptonization. The collapse continues until the rest-mass density exceeds nuclear saturation density ($\rho_c \gtrsim 10^{14}$ g · cm $^{-3}$), leading to core bounce and the formation of a proto-neutron star. Since the collapse and early post-bounce dynamics are consistent with those reported by Longo Micchi et al. (2023) across all models, our focus below is mostly on the post-bounce evolution.

Although all models experience similar collapse and core bounce dynamics, regardless of the initial magnetic field strength, the post-bounce evolution is significantly different. During the collapse, the magnetic fields are amplified primarily through magnetic flux conservation, with the maximum field strength increasing by approximately a factor of $\sim 2 \times 10^3$ at core bounce. The amplification of magnetic fields can be further developed after core-bounce due to magnetic winding and magneto-rotational instability (Duez et al. 2006; Mösta et al. 2015). Fig. 1 shows the magnetization $\beta_{\text{mag}} := P_{\text{gas}}/P_{\text{mag}}$ profiles at different times of different models. For the $B_{\text{pol}} = 10^{12}$ G case, the magnetization β_{mag} is very small even at the early time ($t - t_{\text{bounce}} \lesssim 0.1$ s), and keep decreasing over the simulation. However, although the sigma-magnetization parameter in the polar regions ($\theta \leq 10^\circ$) is increasing and approaches a value of order unity within the simulation’s timeframe, the material in the polar region is only mildly relativistic by the end of the simulation, with $\Gamma_\infty \lesssim 2$ (equation (B15)). Therefore, to discuss the potential properties of the associated GRBs, we focus primarily on the jet energy and the proto-neutron star’s spin-down time.

Jets—To assess the possibility of jet formation, we calculate the energy carried away by the magnetically-dominated ($\beta_{\text{mag}} < 1$) collimated ($\theta \leq 10^\circ$) outflow along z -axis. The right panel of Fig. 3 shows the estimated jet energy for different models. For the $B_{\text{pol}} = 10^{11}$ G case, the jet energy goes beyond 10^{46} erg about 1.25 s after bounce, reaching 10^{48} erg by the end of the simulation. The corresponding isotropic-equivalent energy is: $E_{\text{jet,iso}} = E_{\text{jet}} / [1 - \cos(\theta_{\text{jet}})] \approx 6.42 \times 10^{49}$ erg, while the luminosity is $\sim 9.44 \times 10^{51}$ erg · s $^{-1}$. In contrast, in the $B_{\text{pol}} = 10^{12}$ G case, the jet carries more than 10^{45} erg energy soon after bounce, and approaches 2×10^{50} erg at the end of the simulation. The corresponding isotropic-equivalent energy and luminosity are $\sim 8.03 \times 10^{51}$ erg and $\sim 9.18 \times 10^{54}$ erg · s $^{-1}$, respectively. We note that the jet energy is still increasing at the end of the simulations, and we argue that the $B_{\text{pol}} = 10^{12}$ G case could eventually release comparable amount of jet

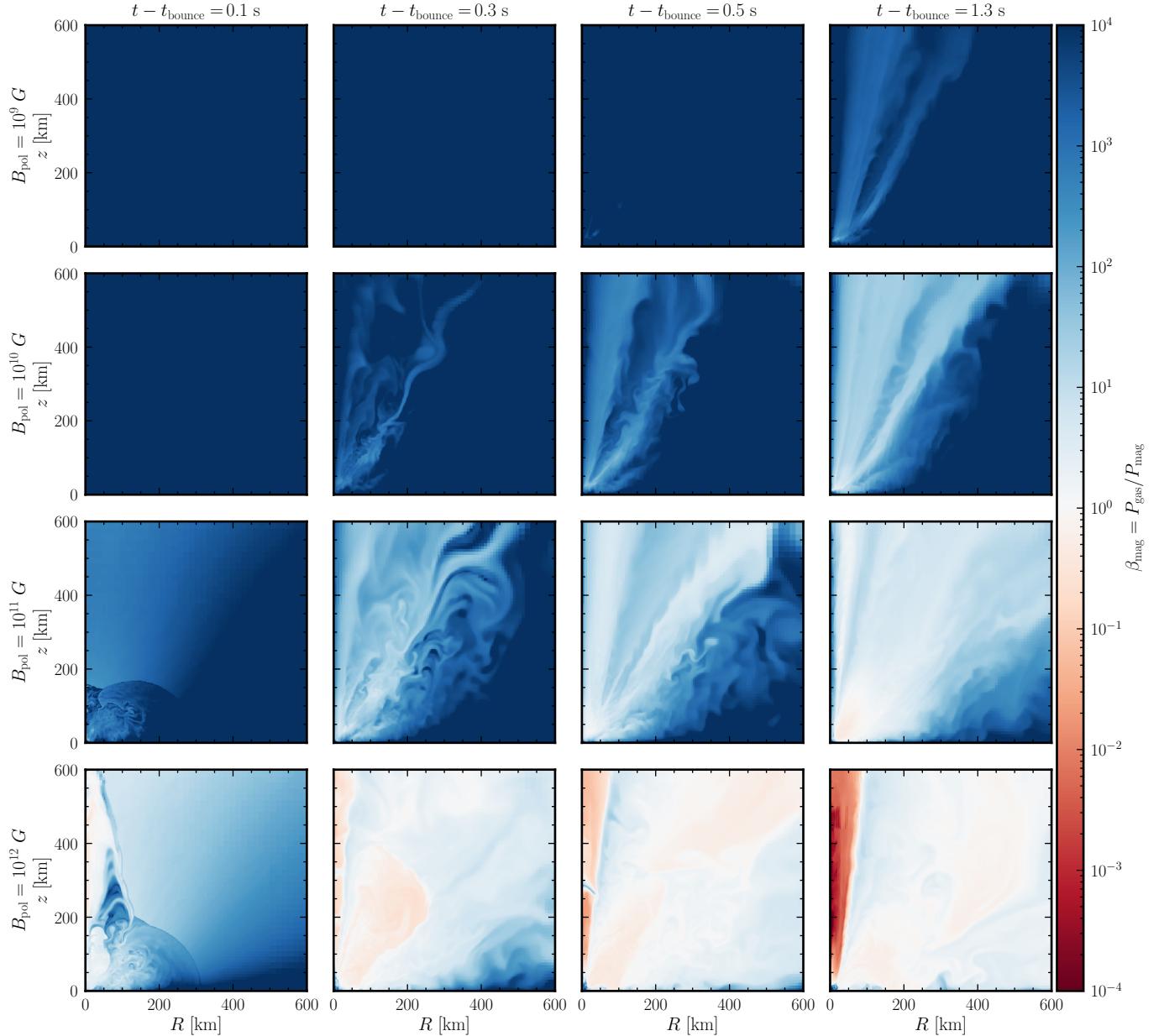


Figure 1. Magnetization $\beta_{\text{mag}} := P_{\text{gas}}/P_{\text{mag}}$ profiles (colour map) at different times (from left to right columns) of different models (from top to bottom rows). The blue regions show the matter-dominating regions ($\beta_{\text{mag}} > 1$) while the red regions represent the magnetic-dominating ones ($\beta_{\text{mag}} < 1$). For the $B_{\text{pol}} = 10^{12} \text{ G}$ case, significant parts of the shown regions are magnetic-dominated, and a strongly magnetized jet-like structure is formed at the polar region. The corresponding magnetization β_{mag} can decrease below 10^{-4} .

energy as in GRB 230307A, whose isotropic-equivalent energy is $\gtrsim 4 \times 10^{52} \text{ erg}$ (see, e.g., [Du et al. \(2024\)](#); [Rastinejad et al. \(2024\)](#)).

The jet is powered by the spin-down of the magnetized proto-neutron star. The duration of the spin-down phase is related to the duration of the gamma-ray burst emission ([Metzger et al. 2008](#)). We estimate the spin-up/down time via $\tau := J_{\text{PNS}}/\dot{J}_{\text{PNS}}$, where J_{PNS} is the angular momentum of the proto-neutron star and \dot{J}_{PNS} is its time derivative. In all our simulations, the proto-

neutron star spins up during the accretion phase, and transitions to a spin-down phase at later times (see e.g. Fig. 2). The stronger the magnetic fields, the sooner the star starts to spin down. At the end of the simulations, the spin-down times are $\{578, 53, 14, 11\} \text{ s}$ in the $B_{\text{pol}} = \{10^9, 10^{10}, 10^{11}, 10^{12}\} \text{ G}$ cases, respectively. For comparisons, the duration T_{90} (time interval comprising 90% of the energy emitted in the 50-300 keV band) of the bursts of GRB 211211A and GRB 230307A are about 34 s ([Mei et al. 2022](#)) and 33 s ([Dichiara et al. 2023](#)) ac-

cording to Fermi observatory, respectively. As the spin-down timescale is mostly affected by the mass ejection, which decrease over time, our values for τ should be regarded as lower limits an for that we consider our findings for spin-down time of the $B_{\text{pol}} = \{10^{11}, 10^{12}\}$ G cases compatible to the duration of the two GRBs of interest.

Matter outflows—In addition to the jet, uncollimated outflows are observed in the simulations. The left and middle panels of Fig. 3 show a comparison of the total ejected mass and energy across different models. By the end of the simulations, the ejected mass and energy in the $B_{\text{pol}} = 10^{12}$ G are roughly an order of magnitude higher than in the $B_{\text{pol}} = 10^{11}$ G case and two orders of magnitude higher than in the $B_{\text{pol}} \leq 10^{10}$ G cases. Our estimated explosion energy is comparable to [Dessart et al. \(2007\)](#).

Magnetic fields play a crucial role in shaping matter outflows, as they can accelerate material and affect whether it is ejected before or after reaching weak equilibrium. Figs. 4 and 5 illustrate the impact of different magnetic field strengths on the ejecta properties. In the weakly magnetized cases ($B_{\text{pol}} \leq 10^{10}$ G), the ejecta is predominantly neutron-poor ($Y_e \approx 0.5$) and slow ($v_\infty \lesssim 0.2c$). For $B_{\text{pol}} = 10^{11}$ G, while the early evolution is similar to the $B_{\text{pol}} \leq 10^{10}$ G models, the material remains significantly more neutron-rich at later times, with electron fractions ranging from 0.25 to 0.5. In the strongly magnetized case ($B_{\text{pol}} = 10^{12}$ G), the outflow properties differ markedly, with a large amount of extremely neutron-rich material ($\sim 0.133 M_\odot$ for $Y_e \leq 0.25Y_e$), as well as higher entropy and faster-moving ejecta. At the end of the simulation, the protoneutron star continues to rotate rapidly, retaining a significant amount of rotational energy ($\sim 10^{52}$ erg), without collapsing. However, it remains unclear how, or if, this energy will be transferred to the ejecta or other channels.

Kilonova emission—Given the significant ejection of material with $Y_e \lesssim 0.25$, the MHD-driven winds of our rotating WD provide an interesting site for r -process nucleosynthesis ([Burbidge et al. 1957](#); [Lippuner & Roberts 2015](#); [Siegel et al. 2019](#); see [Perego et al. 2022](#) and [Cowan et al. 2021](#) for recent reviews), potentially producing a broad range of nuclei. Ejecta with low electron fractions are linked to the production of lanthanides and actinides, that can significantly increase the opacity of the ejecta, shifting the emission towards infrared wavelengths ([Barnes & Kasen 2013](#); [Tanaka & Hotokezaka 2013](#); [Lippuner & Roberts 2015](#)). In contrast, ejecta with higher Y_e values ($Y_e \gtrsim 0.3$) produce low-opacity material, which dominates the blue part of the kilonova spectrum, further influenced by neutrino interactions from the central remnant.

In Fig. 6, we present the results of our kilonova modeling in the AB magnitude system using the J and Ks

Model	$M_{\text{ej}} [M_\odot]$	$E_{\text{ej}} [\text{erg}]$	$E_{\text{jet}} [\text{erg}]$	Spin-down time [s]
$B_{\text{pol}} = 10^9$ G	3.3×10^{-3}	1.0×10^{49}	—	578
$B_{\text{pol}} = 10^{10}$ G	3.6×10^{-3}	9.9×10^{48}	—	53
$B_{\text{pol}} = 10^{11}$ G	2.7×10^{-2}	1.05×10^{50}	1.02×10^{48}	14
$B_{\text{pol}} = 10^{12}$ G	2.6×10^{-1}	1.13×10^{51}	1.67×10^{50}	11

Table 1. Summary of simulation results. Ejecta mass, energy and jet energy are measured at the end of the simulations, corresponding to 1.53 s, 1.76 s, 1.98 s, and 1.3 s post-bounce time, respectively.

Gemini filters, where the grey-radiation transport code SNEC is expected to perform better at later times, when the black-body spectrum peaks in the infrared. The showed bands take into account uncertainties in ejecta opacity, velocity, and total mass. We vary the maximum opacity between $\kappa = 5 \text{ cm}^2 \cdot \text{g}^{-1}$ and $\kappa = 25 \text{ cm}^2 \cdot \text{g}^{-1}$, using the κ fit as a function of Y_e from [Wu et al. \(2022\)](#). Given the dynamic nature of the system at the time of data extraction, we account for uncertainties in the radial velocity profile by considering both the instantaneous and asymptotic velocities of the ejecta fluid elements. Since our simulations are time-limited, not all unbound material has yet reached the extraction radius, and additional ejection could occur beyond the simulation’s timeframe. To account for this, we consider the total ejecta mass M_{ej} between the amount that had exited the grid by the end of the simulation ($\sim 1.6 \times 10^{-1} M_\odot$) and a possible upper limit of $\sim 3 \times 10^{-1} M_\odot$, reflecting potential later ejections. For reference, the total unbound mass at the end of the simulation is $2.6 \times 10^{-1} M_\odot$ (see Table 1). Superimposed over the bands are the afterglow-subtracted observations of GRB 211211A ([Rastinejad et al. 2022](#)) and GRB 230307A ([Fermi GBM Team 2023](#)), taken from table 4 in [Rastinejad et al. \(2024\)](#). Overall, our kilonova modeling shows reasonable agreement with the data, following the predicted reddening of the electromagnetic emission over time. We note that different orientations of the system relative to the observer would further affect the AB magnitudes due to the strong dependence of the Y_e value on the position and polar angle of the ejecta’s fluid elements. However, for simplicity, we do not account for viewing angle effects in this analysis.

4. DISCUSSIONS AND CONCLUSIONS

In this work, we demonstrated that the accretion-induced collapse of strongly magnetized, rapidly rotating white dwarfs are very promising engines of long gamma-ray bursts associated with kilonova emissions. We present second-long axisymmetric GR ν MHD simulations of AIC, showing that magnetic fields play a critical role in jet formation and non-relativistic mass ejection, as summarised in table 1. Moreover, we model the kilonova emission by evolving the initial ejecta profiles via 1D radiation hydrodynamics simulations with

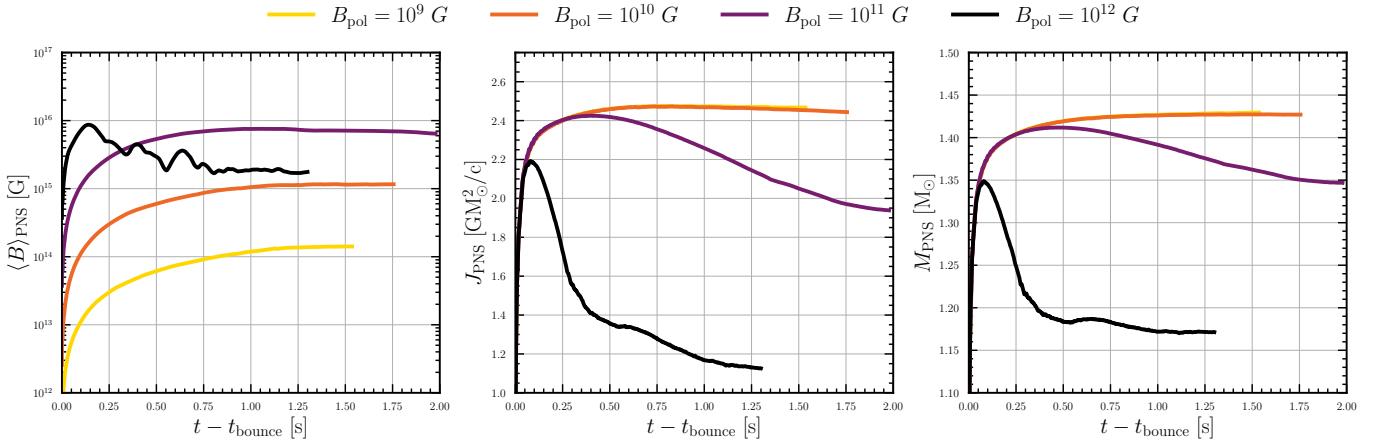


Figure 2. Evolution of the rest-mass density averaged magnetic fields $\langle B \rangle_{\text{PNS}}$ (left panel), angular momentum J_{PNS} (middle panel), and rest-mass M_{PNS} (right panel) of the proto-neutron stars.

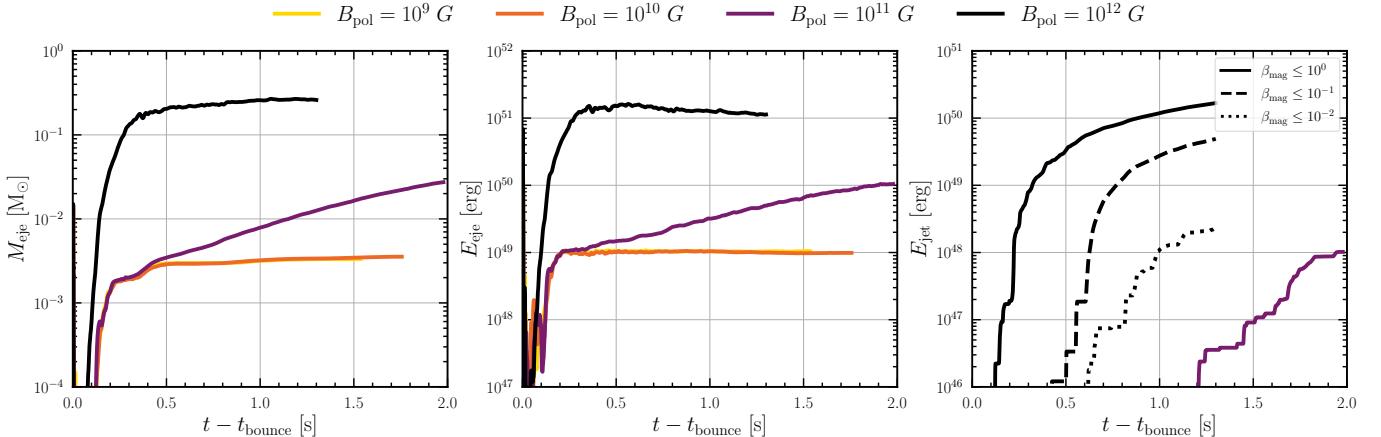


Figure 3. The total ejected mass and energy (left and middle panels), and the jet energy (right panel) of accretion-induced collapse WDs as functions of time. Strong magnetically driven winds enhance the total amount of mass and energies. The jet energy counts the energy carried away via magnetically-dominated collimated outflow that is leaving along the z direction in the region where $\theta \leq 10^\circ$. Solid, dashed, and dotted lines show the same diagnostic energy and with extra constraints on the magnetization: $\beta_{\text{mag}} \leq 1, 0.1, 0.01$, respectively.

realistic opacity and heating rates from radioactive decay. The AB magnitudes in the Gemini J and Ks bands from our kilonova model, when adjusted for uncertainties in ejecta velocity and total mass, agree well with the observations of GRB 211211A and GRB 230307A. The evolution of the emission follows the expected reddening over time, as neutron-rich material leads to infrared-dominated light curves.

Our simulations reveal that highly collimated, magnetically dominated outflows are possible in AICs. In the strongest magnetic field case ($B_{\text{pol}} = 10^{12}$ G), a jet-like funnel forms, with the magnetization parameter (β_{mag}) dropping below 10^{-4} . By the end of the simulation, the jet energy surpasses 10^{50} erg and continues to increase, with an estimated proto-magnetar spin-down time of $\mathcal{O}(10)$ s. These findings align with the proto-

magnetar spin-down model of Metzger et al. (2008), supporting the idea that AICs of strongly magnetized WDs can power GRBs.

Magnetically driven winds significantly enhance mass ejection and alter its composition. In the $B_{\text{pol}} = 10^{12}$ G case, much of the ejecta is neutron-rich, with electron fractions (Y_e) below 0.25, and some as low as 0.1. This material has the potential for full r -process nucleosynthesis, contributing to heavy element production and kilonova emissions.

While only the $B_{\text{pol}} = 10^{12}$ G case produces a powerful jet energy and neutron-rich ejecta, the $B_{\text{pol}} = 10^{11}$ G case may also power GRBs and kilonova over longer timescales than those we have simulated. The mass ejection and jet energy continue to increase by the end of the simulation, and the ejecta becomes progressively more

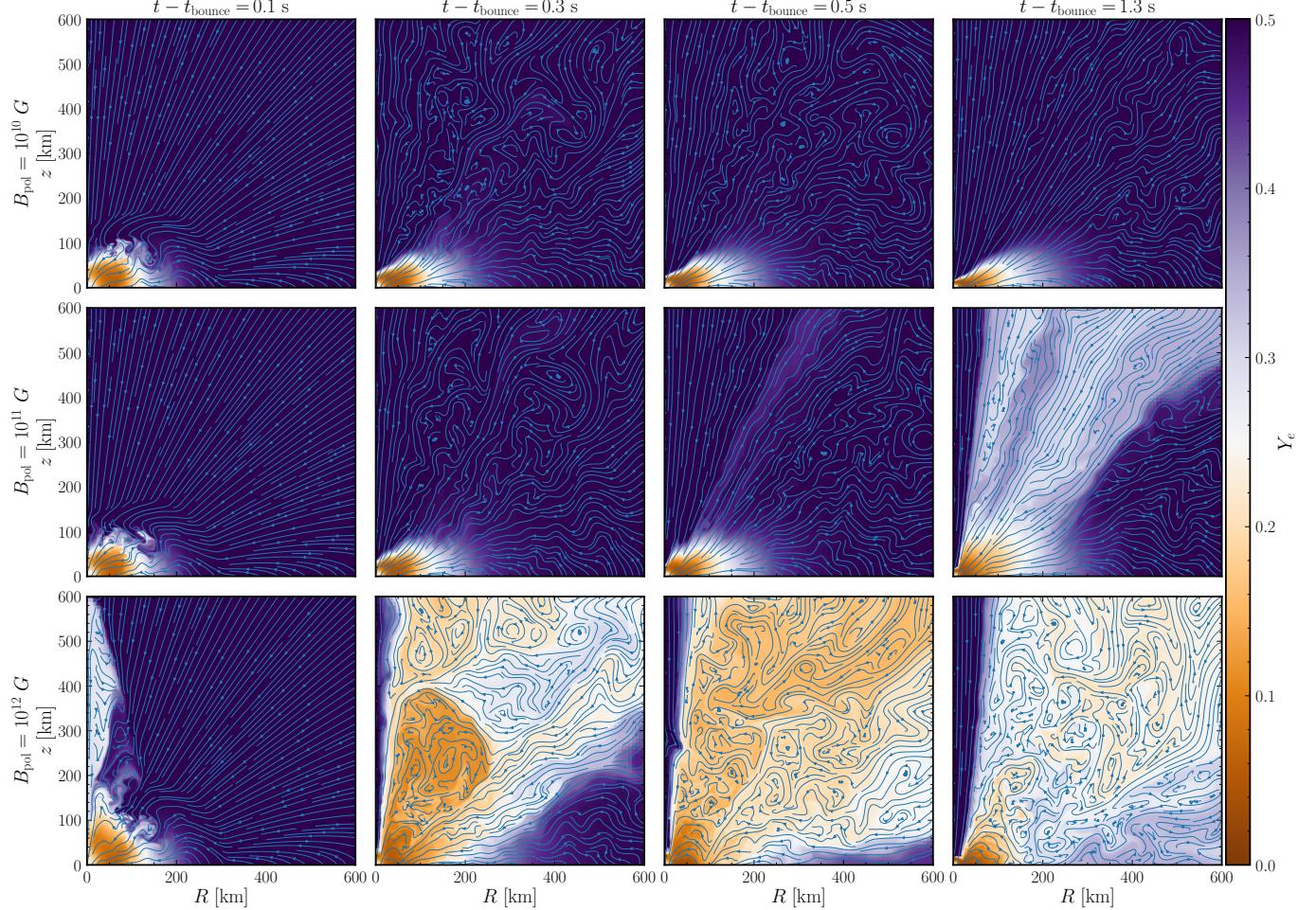


Figure 4. Electron fraction Y_e profiles (colour map) and the magnetic field lines (blue streamlines) at different times (from left to right columns) for different models (from top to bottom rows). The $B_{\text{pol}} = 10^9$ G case is not shown, as it looks identical to the $B_{\text{pol}} = 10^{10}$ G case. Neutron-rich outflow is found when the materials are accelerated out due to strong magnetically driven wind before reaching β -equilibrium. As shown in the figure, when the magnetic field is sufficiently strong (see also Fig. 1), low Y_e materials can be accelerated and escape the disk region. It is clearly visible that the pattern of the electron fraction profile follows the magnetic field lines.

neutron-rich. With further magnetic field amplification, this case may eventually launch jets and eject considerable amounts of neutron-rich material.

We conclude that AIC of strongly magnetized, rapidly rotating WDs offers a natural explanation for events like GRB 230307A and GRB 211211A. In the strongest magnetized model, the jet energy and proto-magnetar spin-down timescale align with long GRBs, and the dim explosion combined with neutron-rich ejecta matches the observed light curves. Thus, AICs are strong candidates for the progenitors of long GRBs with kilonova emissions.

Three-dimensional simulations with higher resolution and longer durations are necessary to further establish the observational implications of our models for LGRBs accompanied by kilonovae. To account for the lack of dynamo amplification and other effects in 2D, we have adopted high initial magnetic fields and rapid rotation,

as discussed in Sec. 2. Future work will address these limitations through longer-term and more detailed 3D simulations, which will enable us to accurately capture magnetic field amplification, jet dynamics, and the role of instabilities such as the kink instability, which are not fully resolved in axisymmetric simulations.

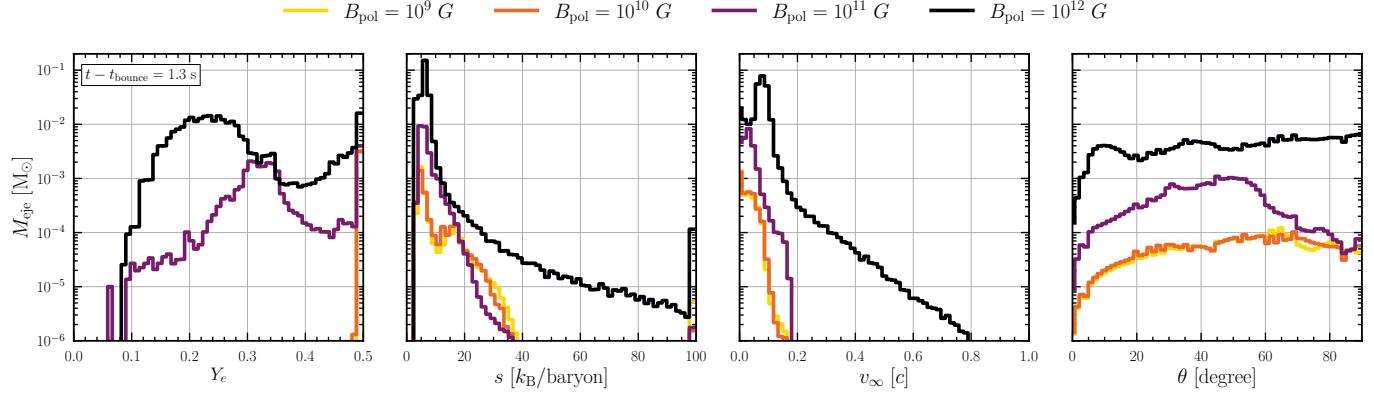


Figure 5. 1-D histograms of the ejecta of an accretion-induced collapse WD up to 1.3 s after bounce. These plots compare the distribution of the ejected mass as functions of electron fraction Y_e , entropy per baryon s , asymptotic velocity v_∞ , and angle ($\theta \in [0^\circ, 90^\circ]$ from pole to equatorial plane) with different initial magnetic fields. The stronger the magnetic fields, the more neutron-rich the ejecta are. Notably, there is a significant amount of very neutron-rich ejecta (i.e. $0.133 M_\odot$ for $Y_e \leq 0.25$) in the case of $B_{\text{pol}} = 10^{12} \text{ G}$.

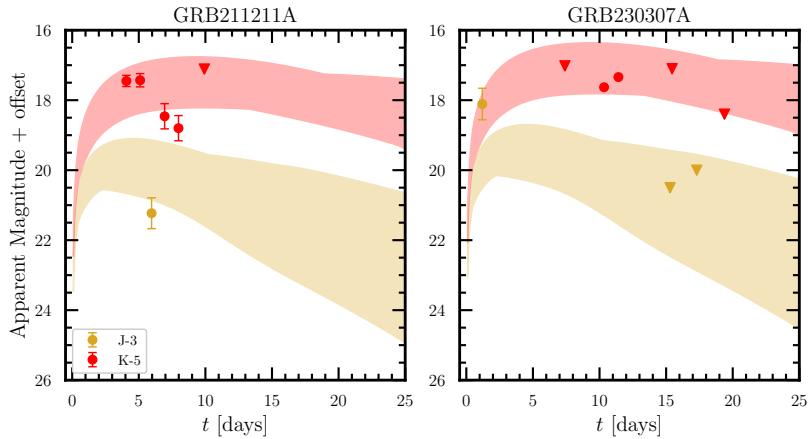


Figure 6. GRB 211211A and GRB 230307A data (dots with error bars and upper limits marked by triangles) compared with SNEC AB apparent magnitudes in Gemini J and Ks filters, with offsets applied for readability. The displayed data correspond to the model with $B_{\text{pol}} = 10^{12} \text{ G}$ (see table 1 for ejecta properties). The bands account for uncertainties in ejecta velocity, total mass, and maximum opacity. We find that our model, with the strongest magnetic field and neutron-rich ejecta, is in good agreement with the observational data. Assumed distances are 350 Mpc for GRB 211211A and 291 Mpc for GRB 230307A.

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Software: We have modified RNS (Stergioulas & Friedman 1995) to generate the initial data. The dynamical simulations of this work were produced by utilising Gmunu (Cheong et al. 2020, 2021, 2022, 2023; Ng et al. 2024). Neutrino rates are provided by the WeakRates module in WhiskyTHC (Radice et al. 2022). The ejecta profile is evolved using the 1D radiation hydrodynamics code SNEC (Morozova et al. 2015; Wu et al. 2022). The data of the simulations were post-processed and visualised with yt (Turk et al. 2011), NumPy (Harris et al. 2020), pandas (pandas development team 2020; Wes McKinney 2010), SciPy (Virtanen et al. 2020) and Matplotlib (Hunter 2007; Caswell et al. 2023).

APPENDIX

A. LOW DENSITY TREATMENT

The rest-mass density of ejecta can be lower than the lowest allowed rest-mass density of a given nuclear equation-of-state table. To capture the dynamics of the low density gas, we extend the equation-of-state by assuming the gas as a mixture of electrons, ions, and photons. In particular, the pressure of the low density gas is given as the sum of ideal-gas and radiation:

$$\begin{aligned} P_{\text{low}} &= P_{\text{gas}} + P_{\text{rad}} \\ &= (\Gamma - 1) \rho \beta_v T + a_{\text{rad}} T^4 / 3, \end{aligned} \quad (\text{A1})$$

where Γ is the adiabatic index, which is chosen to be 5/3. $a_{\text{rad}} = 4\sigma_{\text{SB}}/c$ is the radiation constant, and σ_{SB} the Stefan-Boltzmann constant. The inverse specific heat β_v is a function of T and Y_e obtained by matching the pressure to the table at the lowest density point (i.e. $P_{\text{low}}(\rho_{\text{min}}, T, Y_e) = P_{\text{tab}}(\rho_{\text{min}}, T, Y_e)$):

$$\beta_v(T, Y_e) = \frac{P_{\text{tab}}(\rho_{\text{min}}, T, Y_e) - P_{\text{rad}}(T)}{(\Gamma - 1)\rho_{\text{min}} T}. \quad (\text{A2})$$

The specific internal energy ε and the sound speed c_s needed for the evolution, are given by

$$\varepsilon = \varepsilon_{\text{gas}} + \varepsilon_{\text{rad}} = \beta_v T + a_{\text{rad}} T^4 / \rho, \quad (\text{A3})$$

$$c_s^2 = (c_s^{\text{gas}})^2 + (c_s^{\text{rad}})^2 = \frac{\Gamma P_{\text{gas}}}{\rho + \Gamma \rho \varepsilon_{\text{gas}}} + \frac{1}{3}. \quad (\text{A4})$$

The specific entropy s is needed for post-processing, which is given by

$$\begin{aligned} s &= s_{\text{gas}} + s_{\text{rad}}, \\ &= \left\{ \ln \left[\left(\frac{m_u}{\rho} \right) \left(\frac{2\pi m_u k_B T}{h^2} \right)^{3/2} \right] + \frac{5}{2} \right\} \\ &\quad + \frac{4}{3} a_{\text{rad}} T^3 \left(\frac{m_u}{\rho} \right) \left(\frac{1}{k_B} \right), \end{aligned} \quad (\text{A5})$$

where k_B is the Boltzmann constant, h is the Planck constant, and m_u is the atomic unit.

In the conserved to primitive conversion, we often need to obtain temperature T from a given specific internal energy ε . In this case, the temperature can be obtained

by solving

$$f(T) := 1 - \frac{\beta_v T + a_{\text{rad}} T^4 / \rho}{\varepsilon} = 0. \quad (\text{A6})$$

In practice, we solve this equation via the Newton-Raphson method.

In this work, the floor rest-mass density is set to be 3 orders of magnitude smaller than the lowest available of the equation-of-state table.

B. DIAGNOSTICS

The matter is identified as unbound when it fulfills the Bernoulli criteria and is outgoing. In particular, we locate the unbound matter everywhere in the computational domain by checking

$$f_{\text{ub}} = \begin{cases} 1, & \text{if } h_{\text{tot}} u_t \geq -h_{\min} \text{ and } v_r > 0 \\ 0, & \text{otherwise} \end{cases}, \quad (\text{B7})$$

where $h_{\text{tot}} = 1 + \varepsilon + P/\rho + b^2/\rho$ is specific enthalpy with magnetic field contribution, b^2 is the contraction of the magnetic fields in the fluid frame, h_{\min} is its minimum allowed values for a given equation-of-state, $u_t = W(-\alpha + \beta_i v^i)$, and v_r is the radial velocity. This way, $f_{\text{ub}} = 1$ when the fluid at a point is unbound while equals to zero elsewhere.

To obtain the histograms, we calculate the ejected mass as follows. The ejected mass per bin is given by

$$\Delta M_{\text{eje}} = \int_0^t dt' \oint_{S_{\text{ext}}} \hat{v}^i D f_{\text{bin}} dA_i + \int_{V_{\text{ext}}} D f_{\text{ub}} f_{\text{bin}} dV, \quad (\text{B8})$$

where $\hat{v}^i = \alpha v^i - \beta^i$, and $D = W\rho$ is the conserved rest-mass density. Here, $f_{\text{bin}} = 1$ in the target bins, while $f_{\text{bin}} = 0$ otherwise. The total ejected mass can be obtained by $M_{\text{eje}} = \sum_{\text{bins}} \Delta M_{\text{eje}}$.

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Similarly, the diagnostic ejecta energy is given by

$$E_{\text{eje}} = \int_0^t dt' \oint_{S_{\text{ext}}} \hat{v}^i \tau_{\text{eje}} dA_i + \int_{V_{\text{ext}}} \tau_{\text{eje}} f_{\text{ub}} dV. \quad (\text{B9})$$

Here, τ_{eje} is the diagnostic ejecta energy density, which is defined as

$$\tau_{\text{eje}} = \epsilon_{\text{int}} + \epsilon_{\text{kin}} + \epsilon_{\text{prs}} + \epsilon_{\text{EM}}, \quad (\text{B10})$$

where ϵ_{int} , ϵ_{kin} , ϵ_{prs} and ϵ_{EM} are the internal, kinetic, pressure contribution, and electromagnetic energy densities. They can be obtained by

$$\epsilon_{\text{int}} = \rho W^2 (\varepsilon - \varepsilon_0), \quad (\text{B11})$$

$$\epsilon_{\text{kin}} = \rho W (W - 1), \quad (\text{B12})$$

$$\epsilon_{\text{prs}} = P (W^2 - 1), \quad (\text{B13})$$

$$\epsilon_{\text{EM}} = B^2 \left(1 - \frac{1}{2W^2} \right) - \frac{1}{2} (B^i v_i)^2, \quad (\text{B14})$$

where ε is the fluid specific internal energy, and ε_0 is a reference zero-point which is obtained by the same rest-mass density and electron fraction but with zero temperature (Betranhandy & O'Connor 2020).

The extraction surfaces S_{ext} of the integration of equations (B8) and (B9) are chosen to be a cylinder with radius $R = 1800$ km and $|z| = 1800$ km, while V_{ext} is the corresponding enclosed region.

The estimated asymptotic Lorentz factor is given by

$$\Gamma_{\infty} = -u_t h_{\text{tot}}, \quad (\text{B15})$$

where we assume that all the thermal and magnetic energy will be converted into the kinetic energy of the fluid. The corresponding asymptotic velocity is: $v_{\infty} = \sqrt{1 - 1/\Gamma_{\infty}^2}$.

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