

Resonant amplification of multimessenger emission in rotating stellar core collapse

Marco Cusinato^{1,*}, Martin Obergaulinger^{1,2,†}, Miguel Á. Aloy^{1,2,‡} and José A. Font^{1,2}

¹Departament d'Astronomia i Astrofísica, *Universitat de València*, Av. Vicent Andrés Estellés 19, 46100 Burjassot (València), Spain

²Observatori Astronòmic, *Universitat de València*, 46980 Paterna, Spain



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In a series of axisymmetric core-collapse supernova simulations extending up to ~ 2 s, we identify a regime of precollapse central rotation rates (~ 1 Hz) that greatly enhances the emission of gravitational waves (GWs) during extended periods of time after bounce. The enhancement is a consequence of the resonance between the frequency of the fundamental quadrupolar 2f mode of oscillation of the protoneutron star and the frequency of the epicyclic oscillations at the boundary of the inner core. We observe periods of about several hundred milliseconds each where the resonance is active. The GW emission enhancement produces a correlated resonant modulation of the associated neutrino signal at the same frequencies. With GW frequencies of $\mathcal{O}(1$ kHz) and strain amplitudes within the sensitivity curves of current and next-generation interferometers at distances of $\mathcal{O}(1$ Mpc), this resonant-amplification mechanism may represent a potential game changer for unveiling the supernova explosion mechanism through multimessenger astronomy.

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I. INTRODUCTION

Core-collapse supernovae (CCSNe) mark the final stage in the evolution of massive stars ($M_{\text{ZAMS}} \gtrsim 8 M_{\odot}$) and are among the prime sources of multimessenger emission as they release gravitational waves (GWs), neutrinos, and electromagnetic signals. The detection of 25 neutrinos from Supernova 1987A [1–3], along with its electromagnetic counterpart [4,5], demonstrated that multimessenger observations of CCSNe are feasible.

Modern neutrino observatories, such as Super-Kamiokande [6], IceCube [7], and KM3NeT [8], are expected to detect tens of thousands of neutrinos from a Galactic CCSN [9] or even the collective contribution of cosmological supernovae to the diffuse supernova neutrino background [10]. On the GW side, advanced detectors like Advanced LIGO [11], Advanced Virgo [12], and KAGRA [13] are sensitive to CCSNe within several kiloparsecs. Despite dedicated efforts by the LIGO-Virgo-KAGRA collaboration to detect CCSNe at distances of up to 30 Mpc, no signals have been recorded [14,15]. Nonetheless, prospects for detecting extragalactic events remain promising. Future observatories such as Einstein Telescope (ET) [16] and Cosmic Explorer (CE) [17] are expected to substantially extend the detection range, while emerging insights into processes that enhance

the amplitude of the GW emission from CCSNe may further improve detection chances.

GWs and neutrinos offer complementary insights into CCSN dynamics. For instance, neutrinos were crucial in estimating the gravitational binding energy released in Supernova 1987A [18]. In contrast, GWs reveal signatures of multidimensional fluid instabilities—including protoneutron star (PNS) convection [19,20], stationary accretion shock instability (SASI) [21–23], and prompt convection—emerging during the early stages of the explosion. Moreover, they can constrain the nuclear equation of state (EOS) [24,25] and yield estimates of key PNS parameters [26–28].

The rotation of the progenitor star is pivotal in shaping the waveform of the resulting GW signal and, consequently, the detectability of the event [29,30]. Specifically, when the core of the progenitor star achieves a certain rotational speed, resonances can develop between the PNS oscillation modes and the core's rotational frequency. Such resonances have been identified in simplified two-dimensional (2D) models of rotating relativistic stars [31] and in the collapse of neutron stars to strange quark stars triggered by phase transitions [32]. Additionally, during the early phases ($\lesssim 100$ ms post-bounce) of CCSNe simulations, rotation can excite specific PNS modes [33]. However, the full evolution of these resonances through the entire operation of the CCSNe engine remains to be comprehensively explored. Previous studies indicate that particular precollapse rotation rates can notably amplify the GW amplitude, introduce oscillatory modes, and imprint corresponding signatures in the neutrino signal [33].

In this work, we delve into the rich tapestry of multimessenger signals emerging from CCSN simulations with different initial rotation rates over a span of ~ 2 s. Our findings reveal that when the stellar core rotates at an intermediate rate (≈ 1 Hz), it triggers resonant modes in the PNS that imprint synchronized modulations on the neutrino signal.

*Contact author: marco.cusinato@uv.es

†Contact author: martin.obergaulinger@uv.es

‡Contact author: miguel.a.aloy@uv.es

Remarkably, these resonant GWs generate substantial strains, even surpassing those at bounce, and manifest within mere hundreds of milliseconds after collapse. This striking discovery not only deepens our understanding of the dynamic interplay in stellar explosions but also paves the way for insights into their multimessenger signatures.

The remainder of this paper is organized as follows: Sec. II outlines our numerical setup and presents the progenitor model, initial rotation profiles, and nuclear EOS used to perform the CCSN simulations. In Sec. III, we present the main findings of our study. Finally, in Sec. IV we discuss our results and draw our conclusions.

II. METHODS

We performed 2D CCSN simulations using the AENUS-ALCAR code [34–36] that couples special relativistic magnetohydrodynamics with a spectral two-moment neutrino transport scheme and incorporates an approximately relativistic gravitational potential [37]. For densities exceeding $8 \times 10^7 \text{ g/cm}^3$, we employed the SFHo EOS [38], accounting for photons, electrons, positrons, nucleons, and a full ensemble of light and heavy nuclei in nuclear statistical equilibrium. The GW signal was extracted using the quadrupole formula [39]. The simulations were performed in spherical coordinates on grids with 480 logarithmically spaced radial zones and 128 linearly spaced angular zones, while neutrino energies spanning from 0 to 440 MeV were resolved with 12 logarithmically spaced bins. Each run was evolved for 2 s, commencing from the precollapse phase. In Appendix A, we assess the dependence of our results with increased grid resolution.

After an extensive survey of rotating stellar models evolved to several hundreds of milliseconds postcollapse, we identified several simulations that presented resonant coupling between GWs and rotation. From these, we selected one case in which this feature is particularly pronounced. Specifically, it is a Wolf-Rayet star of subsolar metallicity ($Z = 0.02Z_\odot$) and zero-age main-sequence mass of $M_{\text{ZAMS}} = 17M_\odot$, evolved with both rotation and magnetic fields up to the precollapse stage [40]. Its initially high rotational velocity, coupled with enhanced rotational mixing, drives chemically homogeneous evolution. Building on this model, and reconstructing the magnetic field topology from the saturation strengths originally included in the (one-dimensional) stellar evolution model (of the order of 10^9 G inside the iron core), we explored three rotational configurations: the original (SR), with central rotational rate of $\Omega_c = 0.29 \text{ rad/s}$, an intermediate case (IR) with a rate amplified by a factor of 3.5, and a fast-rotating scenario (FR) with a 12-fold increase. Although these increased rotation rates are substantial, they remain well within the computed range for massive stars [41,42].

III. RESULTS

The plus polarization of the GW signal, h_+ , for a source at distance $\mathcal{D}$ (see top panel of Fig. 1) reveals the brief yet intense bounce signal characteristic of rapidly rotating cores [44,45]. In the most rapidly rotating model FR, the signal amplitude reaches $\mathcal{D}h_+ \gtrsim 100 \text{ cm}$. Following the bounce, the

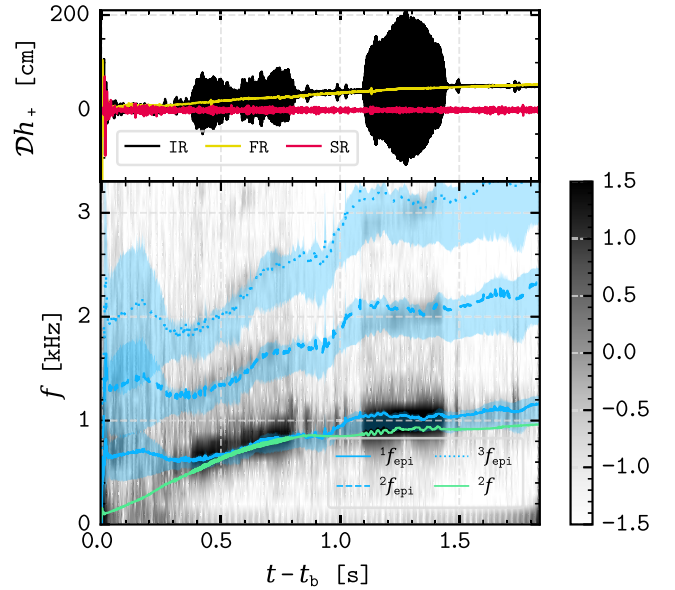


FIG. 1. Top panel: GW strains emitted by a source at distance $\mathcal{D}$ for models SR (red), IR (black), and FR (yellow). Bottom panel: Spectrogram of the GW signal for model IR. Blue lines represent the fundamental frequency (solid), the first (dashed), and the second (dotted) overtones of the epicyclic frequency, with shaded regions showing the associated uncertainties. The green line indicates the fundamental quadrupolar mode, computed using the quasiuniversal relation from Ref. [43].

waveform transitions into a regime of weaker oscillations that remain centered around zero in model SR or gradually drift toward positive values in models IR and FR, which ultimately explode via polar jets.

Unlike models SR and FR, where the oscillatory amplitudes remain relatively constant throughout the simulation, model IR undergoes two distinct phases of pronounced amplitude enhancement, the first occurring between $t \approx 0.4$ and 0.8 s , and the second one between $t \approx 1.1$ and 1.4 s . During these bursts, the strain range widens to ~ 130 and $\sim 300 \text{ cm}$, respectively, reaching levels comparable to the bounce signal observed in model FR.

The bottom panel of Fig. 1 shows the GW spectrogram of model IR, computed using a short-time Fourier transform with a 10 ms time window. The two GW bursts display distinct spectral evolutions: the first rises from $\sim 550 \text{ Hz}$ to $\sim 1 \text{ kHz}$, while the second steadily hovers near 1 kHz . The frequencies and time evolution of the most intense ascending spectral feature align with the predicted fundamental quadrupolar 2f mode (green line computed using as independent variable $x = M_{\text{pns}}/R_{\text{pns}}^3$ instead of the $x = M_{\text{shock}}/R_{\text{shock}}^3$ in the fits provided by Ref. [26]). Moreover, additional power in the spectrogram growing in parallel to the main mode and resembling higher harmonics of the former one appear at twice and thrice these frequencies, albeit with progressively lower amplitudes. Both facts strongly suggest that the rising modes in the spectrogram likely result from PNS f modes. We stress the fact that the green line in Fig. 1 must be regarded as a proxy of the actual 2f -mode frequency. Hence, we refer to this proxy as *approximate* 2f mode. Using the quadrupole formula

including surface terms (see, e.g., Ref. [46]), we pinpoint the origin of the GW strain to the innermost part of the star (hereafter PNS core), which we defined as the region within the isoentropy surface at $4k_B$ /baryon [47], corresponding to a radius of $r_{\text{core}} \sim 20$ km). The pronounced high-amplitude emission arises from a resonant interplay between PNS modes and rotation. The centrifugal force causes fluid elements within the PNS core to oscillate in the direction of the cylindrical radial distance. To quantify these oscillations, we use the epicyclic frequency, defined as

$${}^n f_{\text{epi}} = \frac{n}{2\pi} \sqrt{\frac{2\Omega}{R} \frac{d(\Omega R^2)}{dR}}, \quad (1)$$

where n is the overtone number, R the cylindrical radius, and Ω the angular velocity.

Within the PNS core, we focus on a conical region spanning the colatitude range $[15^\circ, 85^\circ]$. At each radius, we determine the maximum epicyclic frequency and subsequently compute the average and standard deviation of these maxima. This specific colatitude range was chosen to exclude the polar region (prone to numerical artifacts in axial symmetry) and the equator, where converging convective flows compromise the accuracy of the ${}^n f_{\text{epi}}$ evaluation. Recognizing that simply taking the maximum may lead to an overestimate, we conservatively define the lower uncertainty as the largest value between the standard deviation and 10% of the maximum of the epicyclic frequency. The bottom panel of Fig. 1 shows the evolution of these averages (blue lines) and uncertainties (shades) for ${}^1 f_{\text{epi}}$ (solid), ${}^2 f_{\text{epi}}$ (dashed), and ${}^3 f_{\text{epi}}$ (dotted).

Resonance occurs when the fundamental epicyclic frequency intersects or comes sufficiently close to the ${}^2 f$ -mode frequency. To give a first analytic estimate of the latter, we employ the quasiuniversal relations of Ref. [26]. These relations were obtained from nonrotating 1D CCSN simulations. Thus, their approximate formulas may exhibit deviations from our models, yet they provide a rough reference for comparison with the epicyclic frequency.

During the resonance phases, the PNS core oscillates violently over few tens of milliseconds. Since the PNS core and PNS essentially coincide during the second resonant phase, the two terms are used interchangeably in the following discussion. In Fig. 2, we provide a time-sequence visualization of the fluid motion within the PNS core and the associated deformation during the second resonant phase, showing the radial velocity (v_r) during a full oscillation cycle. The PNS core motion takes the form of a large-scale oscillation in which the core undergoes over a full cycle the following dynamical sequence: a roughly isotropic contraction (top left), formation of two vortices, one per hemisphere, with inward/outward radial velocity at the equator/poles (top right), expansion (bottom right), and oppositely oriented vortex pair (bottom left). Over the course of one oscillation, the PNS core changes its quadrupole moment due to the differential between motion in the equatorial and polar regions (second and fourth phase). At the outer border of the PNS (black line), each cycle of the oscillation launches a wave into the postshock region, visible in the four panels as the approximately isotropic patterns of

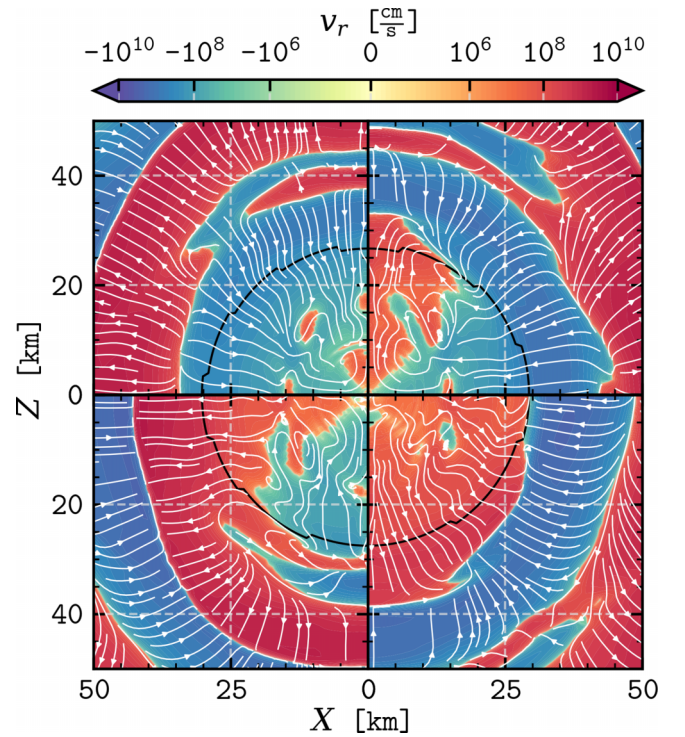


FIG. 2. Snapshots of the radial velocity illustrating a full oscillation cycle of the PNS at (clockwise from the top-left panel) 1.1283, 1.1285, 1.1287, and 1.1291 s. Each of the panels shows one hemisphere only. Black lines represent the PNS radius, defined as the isodensity line at 10^{11} g/cm³. Streamlines represent the velocity field. Background colors indicate inward- (blue) and outward-moving (red) matter.

alternating positive and negative velocities moving away from the PNS.

To further validate the resonance mechanism, we resort to the computed data to give a second estimate of the frequency of the quasiradial ($l = 0$) and quadrupolar ($l = 2$) modes. As in Ref. [31], we use as proxies for the amplitudes of the $l = 0$ and $l = 2$ modes the real parts of the Fourier spectra of the mass density at $r = 25$ km, $\theta = \pi/2$ and of v_θ at $r = 25$ km, $\theta = \pi/4$, respectively. The choice of $r = 25$ km ensures that the extracted signal reflects fluid displacements induced by the oscillations of the PNS core. In Fig. 3, we present the Fourier transform amplitude (FTA) of the density and v_θ , each normalized by their maximum during the second resonant phase (1.15–1.45 s), along with the GW FTA computed in the whole domain ($r \leq 10^{10}$ cm). The GW spectrum (green line) exhibits a series of distinct peaks starting at ~ 1 kHz, spaced at intervals of ~ 1 kHz. These peaks persist up to ~ 7 kHz, though their amplitude steadily diminishes with increasing frequency, a behavior indicative of overtones of the epicyclic frequency.

The spectral peaks of the $l = 0$ mode (red line) and $l = 2$ (blue line) closely align with those of the GW signal and, to a rather good approximation, with the epicyclic frequency and its overtones. The slight shift between the local GW spectral peak spectrum and the epicyclic frequency is likely due to its overestimation. Furthermore, the first peak in the GW spectrum, corresponding to the fundamental epicyclic frequency,

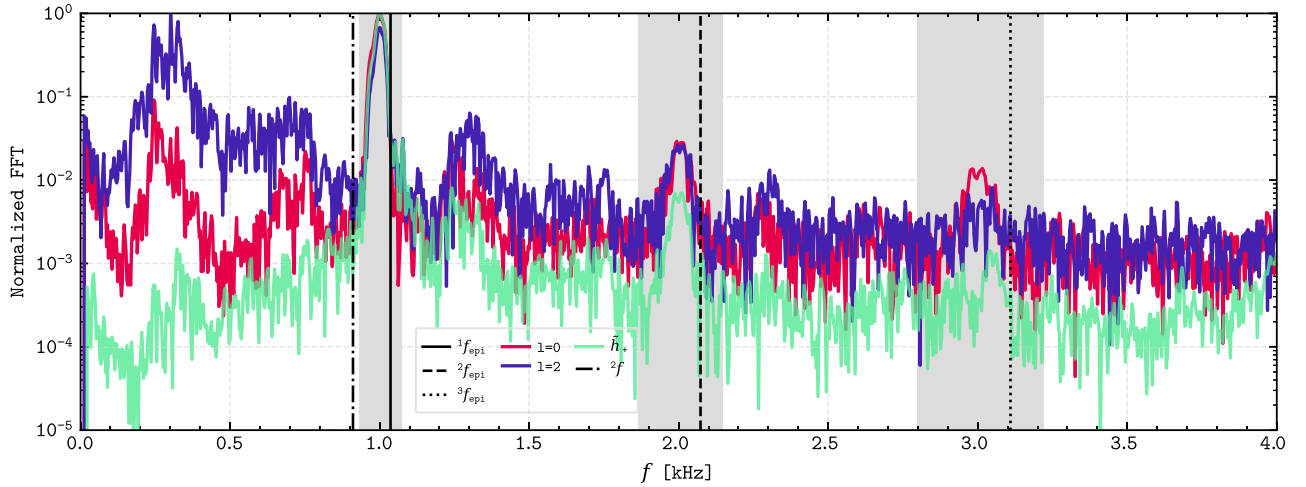


FIG. 3. Normalized Fourier transform of the equatorial density ($l = 0$), v_θ at $\pi/4$ ($l = 2$), both outside of taken the PNS core outer boundary at 25 km, and of the GW signal ($\tilde{h}_+$) for the time interval 1.1–1.4 s. The solid, dashed, and dotted black vertical lines represent the fundamental frequency and the first and second overtones of the average epicyclic frequency, respectively. Shaded regions indicate the uncertainty associated with the frequencies. The dash-dotted line represents the average 2f -mode frequency derived with the relations in Ref. [43].

agrees with the 2f mode estimated via the quasiuniversal relations of Ref. [43]. Together, these findings suggest that the fundamental epicyclic mode is the primary driver of the observed GW emission during the resonant phase.

To assess the detectability of the GW signal, we compute its characteristic strain [48] h_{char} ,

$$h_{\text{char}}(f) = 2f |\tilde{h}(f)|, \quad (2)$$

where $\tilde{h}(f)$ is the Fourier transform of the GW strain $h(t)$ after applying a Hann window $w(t)$,

$$\tilde{h}(f) = \int_{-\infty}^{+\infty} h(t) w(t) e^{-2\pi i f t} dt. \quad (3)$$

The Hann window is defined as

$$w(t) = \frac{1}{2} \left[1 - \cos\left(\frac{2\pi t}{T}\right) \right], \quad 0 \leq t \leq T, \quad (4)$$

and $w(t) = 0$ outside the interval $[0, T]$, where T is the total duration of the signal.

The use of a Hann window reduces spectral leakage associated with the finite time span of the signal, but it also introduces a reduction of the Fourier amplitudes. For a coherent signal, the Hann window has a coherent gain $\langle w \rangle = T^{-1} \int_0^T w(t) dt = 1/2$ [49]. To preserve the correct normalization of the Fourier amplitude, and hence of h_{char} , the windowed Fourier transform is therefore renormalized by the inverse of this coherent gain.

In Fig. 4, we show $h_{\text{char}}/f^{1/2}$ for models SR, IR, and FR at $\mathcal{D} = 1$ Mpc, juxtaposed with the amplitude spectral density (ASD) of current and next-generation interferometers, namely, Advanced LIGO [11,50,51], Advanced Virgo [12,50,51], KAGRA [13,50,51], ET [52,53], and CE [54,55].

Model SR notably exhibits a broad low-frequency hump spanning 40–90 Hz alongside a broad high-frequency peak at 400–1000 Hz. While the low-frequency features fall below the sensitivity range of present and future detectors, the

high-frequency peak marginally surpasses the theoretical sensitivity of CE. Moreover, in the absence of polar jets, no significant power is observed below $f \sim 30$ Hz.

Model FR exhibits a broad plateau below 250 Hz. Its high-frequency peak is shifted down to ~ 400 Hz and is markedly weaker compared to models SR and IR. Contrary to the previous model, a spectral power excess at low frequencies ($\lesssim 30$ Hz), is present, indicating jet formation.

Model IR differs significantly from the others. It displays two low-frequency spectral humps: one at intermediate frequencies (60–200 Hz) and another below 30 Hz. The most striking feature of IR is the unmistakable resonance imprint on its characteristic strain. This model yields the strongest signal, with a peak strain more than twice that observed in

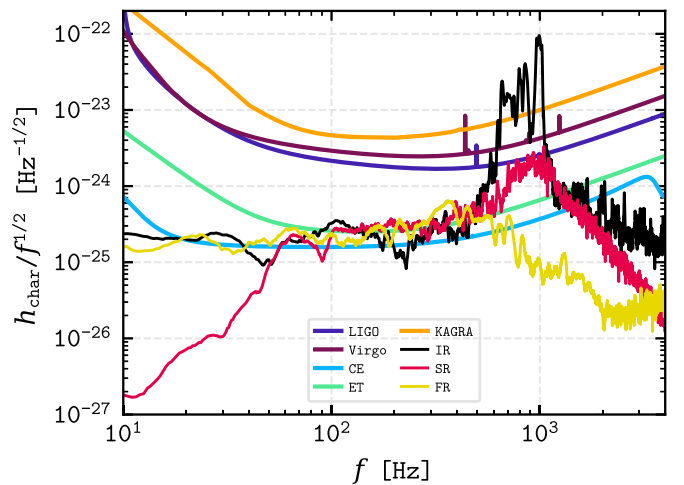


FIG. 4. Characteristic GW spectra for models SR (red), IR (black), and FR (yellow), assuming a source distance of 1 Mpc, compared with the design sensitivity curves of current and next-generation interferometers. h_{char} has been convolved over a window of 10 Hz for better visualization.

SR, translating into a peak ASD more than 10 (100) times larger than in model SR (FR) (see Fig. 4). Four prominent peaks, spanning from ~ 600 Hz to ~ 1 kHz, mark the resonant frequencies during the first and second resonant phases (see Fig. 1). Overtones at $f \gtrsim 2$ kHz and $f \gtrsim 3$ kHz are also present. To estimate the maximum distance at which the GWs from our models remain detectable, we adopt an optimal signal-to-noise ratio (SNR) threshold of 12. This value is higher than the usual threshold of 8 used in recent works [30,56] to account for the intrinsic overestimation resulting from the axisymmetric nature of the simulation. We define the SNR as in Moore *et al.* [48],

$$\text{SNR} = \sqrt{\int_0^{+\infty} d \ln f \left(\frac{h_{\text{char}}(f)}{h_n(f)} \right)^2}, \quad (5)$$

where $h_n(f) = \sqrt{f} \text{ASD}(f)$, in which we used the ASD of the detector noise. The strong resonant features of model IR lie within the detection range of current interferometers such as Advanced LIGO and Virgo beyond the Andromeda Galaxy [57], with a maximum reach of 1 Mpc. Next-generation observatories like CE, with enhanced high-frequency sensitivity, could extend this range to ~ 7.4 Mpc.

In model IR, the neutrino emission shows resonant features at frequencies matching those of the GWs signal, albeit at slightly lower amplitudes. These resonant imprints emerge in the mean energies and luminosities of all three neutrino flavours (ν_e , $\bar{\nu}_e$, and the combined τ and μ neutrinos, ν_x). The two bottom panels of Fig. 5 depict, respectively, the evolution of the global neutrino luminosity (middle panel) and the spectrogram of the electron antineutrino luminosity (bottom panel), both computed with a 10 ms time window, analogous to the analysis conducted for the GW signal.

During the resonance intervals observed in the GW signal (0.4–0.8 and 1.1–1.4 s), the global neutrino luminosity oscillates around its mean value. Although such oscillations are evident in all neutrino flavours, the resonance imprint is most pronounced in the $\bar{\nu}_e$ component, followed by ν_e , and is least apparent in ν_x signal. The spectrogram of the $\bar{\nu}_e$ luminosity (bottom panel of Fig. 5) mirrors that of the GWs, displaying two distinct phases of robust, resonantly excited oscillations at the epicyclic frequency or its overtones (blue lines).

To quantify the correlation between the GW and neutrino signals, we use the matching score [58],

$$\mathcal{M}_{v_i}(t) = \frac{\langle S_{\text{GW}}(f, t) | S_{v_i}(f, t) \rangle}{\sqrt{\langle S_{\text{GW}}(f, t) | S_{\text{GW}}(f, t) \rangle \langle S_{v_i}(f, t) | S_{v_i}(f, t) \rangle}}, \quad (6)$$

where $S_j(f, t)$ ($j = \text{GW}, v_i$) denotes the magnitude of the spectrogram computed with a 10 ms time window and $\langle \cdot | \cdot \rangle$ represents the inner product in frequency space.

The top panel of Fig. 5 depicts the evolution of the matching score between the GWs (bottom panel of Fig. 1) and $\bar{\nu}_e$ spectrograms (bottom panel of Fig. 5; results for the other flavours are similar) for frequencies above 200 Hz, a threshold set to minimize stochastic contributions from convection and SASI. Throughout the simulation, the matching score consistently exceeds 0.6. Notably, during the two resonant phases (0.4–0.8 and 1.1–1.4 s), it approaches unity, unequivocally demonstrating that the amplitude modulations in both signals

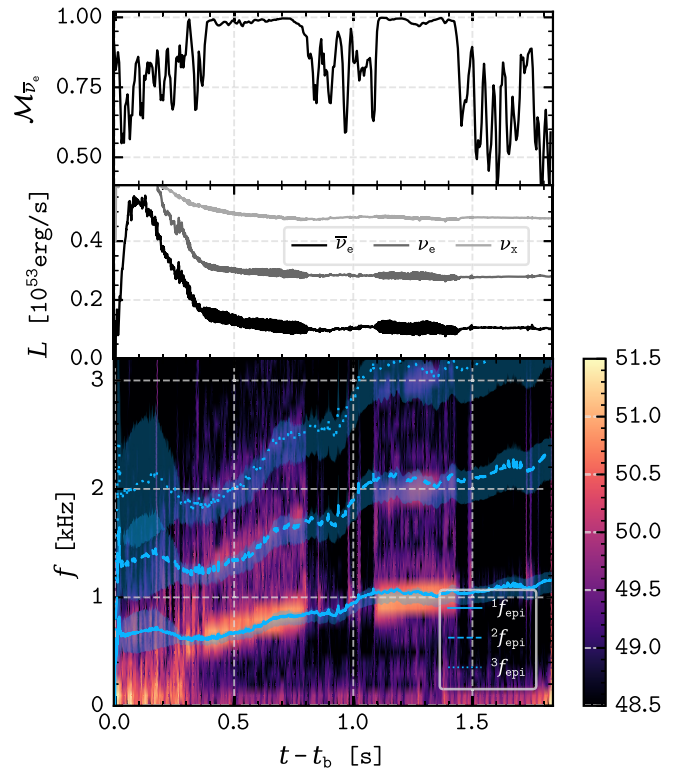


FIG. 5. Top panel: evolution of the matching score between the spectrograms of the $\bar{\nu}_e$ and GW signals. Middle panel: neutrino luminosities of electron antineutrinos (black line), electron neutrinos (dark gray, shifted by 0.2×10^{53} erg/s), heavy lepton neutrinos (light gray, shifted by 0.4×10^{53} erg/s). Bottom panel: spectrogram of the antineutrino luminosities, with the shaded regions and blue lines denoting the same quantities as in Fig. 1. Both panels refer to model IR.

are correlated and originate from the same rotationally driven oscillations of the PNS core.

IV. DISCUSSION AND CONCLUSIONS

We have reported results from 2D CCSN simulations for a model with a prebounce central rotation rate of $\Omega_c \sim 1$ rad/s that reveal extended intervals—lasting several hundred milliseconds—of extraordinarily strong GW emission. Such pronounced signals are absent in otherwise identical progenitors with either slower or faster initial rotation. We attribute this enhanced emission to a resonance between the frequency of the fundamental quadrupolar 2f mode and the frequency of epicyclic oscillations at the outer boundary of the PNS inner core. Notably, while the GW amplitudes are comparable to the bounce signals observed in rapidly rotating cores, they persist for much longer. Crucially, the resonance signatures dominate the most prominent features of the characteristic strain, and their typical frequencies of around 1 kHz renders them compelling targets for ground-based GW detectors. Although higher harmonics are also excited, their detection is impeded by both their high frequencies and lower amplitudes. We also detect imprints of these resonant oscillations in the neutrino emission from the

core, with luminosities and mean energies modulated at the same frequencies, most prominently in the $\bar{\nu}_e$ and ν_e channels.

While our findings suggest an effect relevant for multimessenger observations of CCSNe, further investigation is clearly warranted. In realistic CCSN cores, deviations from axisymmetry arise rapidly; at the rotation rates explored here, nonaxisymmetric instabilities may develop and alter the rotational profile, potentially modifying the conditions for resonance. These effects must be rigorously examined using full three-dimensional (3D) simulations. Moreover, 2D simulations may overestimate GW amplitudes compared to 3D [59], which represents a limitation of the present study. Furthermore, future studies should probe whether this resonance is sensitive to other progenitor properties, such as initial mass or metallicity, and if it can offer constraints on the nuclear EOS.

Bearing in mind these caveats, the pronounced, long-lasting modulation of the GW signal induced by resonance reported in this work could represent a breakthrough for multimessenger detections of CCSNe. Given the substantial GW amplitude enhancement, this resonant signature should remain detectable out to distances of ~ 1 Mpc with current detectors. With CE, the detection horizon extends to ~ 7.4 Mpc. This makes the signal a compelling target for both current and next-generation interferometers.

To provide an estimate of the expected number of detectable events, we approximate the detection rate as

$$\mathcal{R}_{\text{det}} \simeq \mathcal{R}_{\text{CCSN}}(D < 10 \text{ Mpc}) V(d_r) f_{\text{rot}}, \quad (7)$$

where $\mathcal{R}_{\text{CCSN}}(D < 10 \text{ Mpc})$ is the local volumetric core-collapse supernova rate (i.e., within a distance $D < 10 \text{ Mpc}$), $V(d_r)$ is the accessible volume set by the detector horizon distance d_r , and f_{rot} is the fraction of massive-star progenitors that meet the (model-dependent) conditions for the resonance (e.g., sufficiently rapid presupernova core rotation, namely, $\sim 1 \text{ rad/s}$ [60]; see Appendix B). We take $V(d_r) = 4\pi d_r^3/3$ and adopt $d_r \simeq 7.4 \text{ Mpc}$ for CE (as inferred from our SNR analysis). For the local CCSN rate we use the commonly employed empirical estimate of $\sim 1\text{--}2$ CCSN per year within 10 Mpc, with uncertainties dominated by incompleteness due to dust extinction and intrinsically faint events (e.g., Refs. [61–64]). That corresponds to $\mathcal{R}_{\text{CCSN}}(D < 10 \text{ Mpc}) \approx 2.4 \times 10^{-4} \text{ yr}^{-1} \text{ Mpc}^{-3}$. Stellar evolution and population-synthesis studies indicate that only a small fraction of massive-star progenitors are expected to retain presupernova core rotation rates of order $\sim 1 \text{ rad/s}$, with estimates ranging from $f_{\text{rot}} \sim 10^{-2}$ up to $\sim 10^{-1}$ in optimistic scenarios involving chemically homogeneous evolution or strong binary interaction (e.g., Refs. [40,60,65]). As a fiducial, but highly uncertain value one may set $f_{\text{rot}} = 0.03$, yielding

$$\begin{aligned} \mathcal{R}_{\text{det}} &\simeq 0.012 \text{ yr}^{-1} \left(\frac{f_{\text{rot}}}{0.03} \right) \times \left(\frac{d_{\text{hor}}}{7.4 \text{ Mpc}} \right)^3 \\ &\times \left(\frac{\mathcal{R}_{\text{CCSN}}(D < 10 \text{ Mpc})}{2.4 \times 10^{-4} \text{ yr}^{-1} \text{ Mpc}^{-3}} \right), \end{aligned} \quad (8)$$

i.e., of order one event every ~ 80 yr.

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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 A: RESOLUTION STUDY

Two additional high-resolution simulations were performed to assess the dependence of the resonance on spatial resolution. We find that the resonance phenomenon persists, but the associated GW flares become weaker and shorter-lived at higher resolution. This behavior can be explained by small shifts in the f -mode frequency that prevent sustained overlap with the epicyclic frequency, highlighting the strong sensitivity of the resonance to resolution-dependent frequency variations.

All three simulations span the same domain: a radial extent from 0 to 10^{10} cm and an azimuthal angle range of $[0, \pi]$. Model IR uses an innermost grid spacing of 400 m, with $n_r = 480$ radial zones and $n_\theta = 128$ azimuthal zones. Model IR-2 refines the innermost spacing to 300 m, with $n_r \times n_\theta = 602 \times 256$, while model IR-3 maintains 300 m spacing but uses $n_r \times n_\theta = 800 \times 512$.

Models IR-2 and IR-3 exhibit short flares of enhanced amplitude, lasting a few tens of milliseconds and reaching amplitudes of several tens of centimetres, but both the duration and amplitude are reduced compared to IR (Fig. 6, top panel). These flares are more similar in duration and amplitude to models R3 and R4, in which the rotation has been altered (see next section). The bottom panels display the spectrograms. In higher-resolution runs the f -mode frequency, indicated by the darker band of increasing frequency, is shifted upward by $\sim 100 \text{ Hz}$, so the fundamental epicyclic frequency (solid blue lines) and its maximum (blue shading) do not remain aligned with the f mode long enough to drive a strong resonance, producing only short, weak flares.

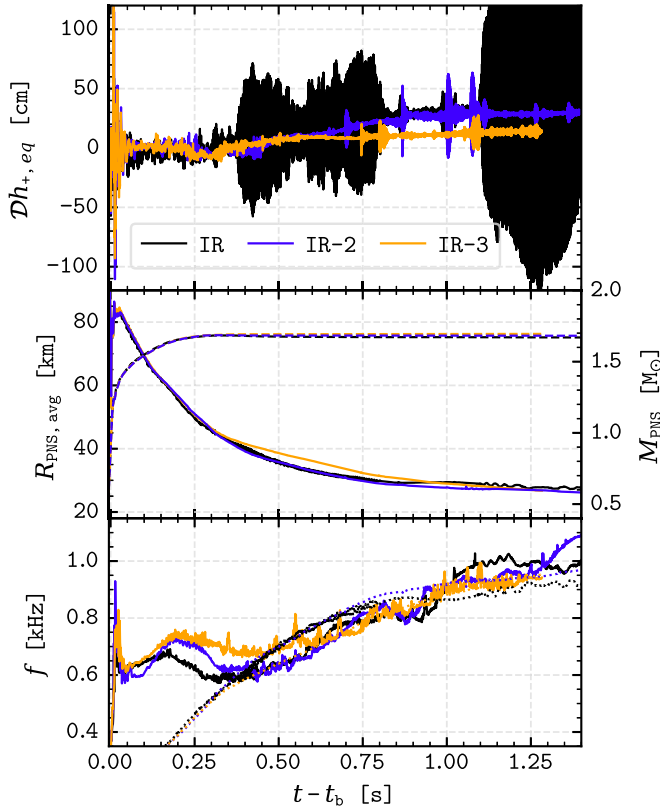


FIG. 6. Top panel: GW strains emitted by a source at distance $\mathcal{D}$ for models IR (black), IR-2 (blue), and IR-3 (orange). Middle panel: evolution of PNS radius (solid lines) and mass (dashed lines). Bottom panel: epicyclic frequency evolution (solid lines) and f mode, computed using Ref. [43] (dotted).

Consistency across resolutions is further confirmed in the bottom two panels of Fig. 6. PNS quantities for IR and IR-2 agree within 5%. Model IR-3 shows a slightly larger and more oblate PNS, due to its faster rotation. The epicyclic frequencies agree within 100 Hz overall, but in the first 0.3 s postbounce, when the initial resonant bursts appear in IR, the higher-resolution models yield slightly higher values. This may delay the crossing with the f mode beyond the computed times or happen in shorter periods, thereby damping the resonance.

As in many other resonant phenomena in physics (e.g., a driven, slightly damped pendulum or the propagation of seismic waves with anelastic attenuation), damping implies that a higher quality factor Q results in a narrower frequency range over which true resonance occurs. The resolution study presented here suggests that the resonance observed has a relatively high Q , so even small detuning between the epicyclic frequency and the PNS f -mode frequency significantly reduces the resonant amplitude. Consequently, time-dependent factors that slightly alter the conditions governing either the f mode or the epicyclic frequency can induce partial quenching or intermittency in the resonant response.

This argument explains how changes in resolution—leading to the mild dynamical differences discussed in previous paragraphs—can cause a detuning between the 2f and f_{epi} modes and, accordingly, a significant reduction in

the GW amplitude. This result should not be interpreted as a lack of physical relevance of the identified resonance. Rather, it suggests that very slight modifications to the initial conditions—specifically, to the central nuclear rotational rate Ω_c —can trigger large-amplitude GWs in these models (as observed for $\Omega_c = 1$ rad/s). The substantial computational cost of higher-resolution models prevents a more detailed exploration to confirm this point. However, this interpretation is further supported by additional simulations based on different presupernova stellar models (not shown in this work) that exhibit similar resonant behavior.

We also attribute to the high Q the partial quenching of the resonance observed between the two main flares in model IR. The low-activity period coincides with a phase during which extended equatorial inflows strike the PNS surface. Most of the high-specific-angular-momentum material impacting the PNS is not assimilated (as indicated by the absence of a clear increase in PNS mass); instead, it rebounds and is eventually expelled through the polar regions into the collimated supernova ejecta. We tentatively identify the random variations induced by these inflows on the PNS as a plausible source of detuning between the f mode and epicyclic frequencies.

In summary, the resonance is robust to changes in resolution, but its GW signature depends sensitively on the relative values of the f mode and epicyclic frequencies. Further high-resolution simulations are required to quantify the enhancement produced by this resonance and to constrain the parameter space in which it arises.

APPENDIX B: ROTATION DEPENDENCE

Two additional simulations exploring the rotation parameter space were carried out to test how the resonance depends on the initial rotation rate. These correspond to model R3, initialized with three times the rotation rate of the progenitor, and model R4, with a fourfold increase. These models are configured with nuclear rotational rates near the value at which we observed resonant behavior, specifically $\Omega_c = 0.87$, 1, and 1.16 rad/s. The results show that the resonance is highly sensitive to rotation: increasing or decreasing the initial rotation changes both the timing and the character of the GW signal.

The top panel of Fig. 7 compares the GW strains of models R3, R4, and IR. The reference model IR exhibits a continuous, sustained resonance over 0.4–0.8 s. Model R3, instead, shows a strong but short-lived enhancement lasting about 100 ms, beginning at 0.35 s, followed by a quiescent phase until 0.9 s. Model R4 produces yet another pattern: a sequence of intermittent flares of elevated GW amplitude, each lasting only a few tens of milliseconds. Together, these comparisons demonstrate that the resonance peaks for a specific rotation, while weakens or fragments for increased or decreased spin.

The middle panel of Fig. 7 shows the PNS mass (dashed lines) and radius (solid lines). The three models produce very similar PNS structures, with only a small mass offset: Model R3 is more massive by about $\sim 0.05 M_\odot$, while the radii remain nearly identical.

The bottom panel of Fig. 7 compares the epicyclic frequency (solid lines) and the ratio of rotational to gravitational energy, $T/|W|$ (dotted lines). In the slower-rotating models

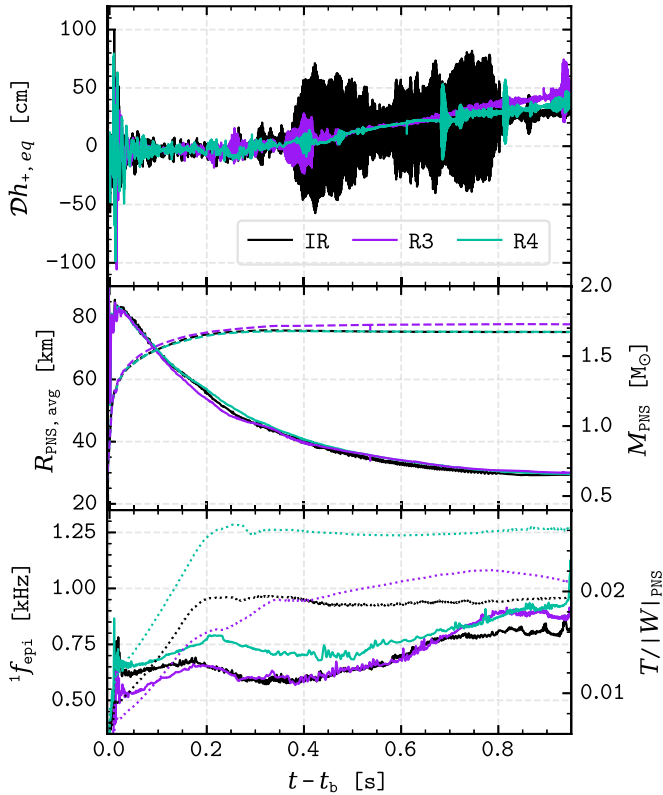


FIG. 7. Same as Fig. 6 but for models IR, R3, and R4, in which the f mode has been replaced with $T/|W|$.

(IR and R3), the epicyclic frequencies are nearly identical between 0.2 and 0.6 s. However, in model R3, $T/|W|$ gradually

increases and eventually surpasses that of IR, even though IR began with faster rotation. After resonance is triggered, the value of $T/|W|$ in IR stabilizes, while it continues to rise in R3. In the fastest-rotating case, R4, both the epicyclic frequency and $T/|W|$ are consistently higher by about ~ 100 and 0.005 Hz, respectively. These differences indicate that changes in the PNS shape and rotation rate can shift the frequency of the f mode, altering the conditions required for the epicyclic frequency to maintain resonance. As a result, sustained enhancements of the GW signal become harder to achieve at higher rotation rates.

The sensitivity to the rotational rate must be interpreted with care. Changes in numerical resolution propitiate a slight detuning between the f mode and epicyclic frequencies, driven by mild differences in the dynamics. We suspect that resolution-dependent effects may lead to conditions in which the resonance is amplified for different initial values of Ω_c . The substantial computational resources required to verify this hypothesis compel us to defer a detailed study to future work.

In summary, the resonance is found within a relatively narrow range around $\Omega_c \approx 1$ rad/s under the idealized conditions of our models, where rotation is imposed on an initially nonrotating stellar structure. Combined with the uncertainties in the actual rotational profiles predicted by stellar evolution models, this makes it very difficult to estimate the fraction of massive, rotating stars in which conditions are favourable for the development of the resonance. Such conditions may occur in, say, $\sim 10\%$ of the tail of faster rotators (i.e., with equatorial projected rotational speeds $v \sin i > 150$ km/s; Ref. [60]). This implies that only about 1% of all massive stars could host favourable conditions for the resonance discussed in this work.

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