

Distinct neutrino signatures and onset condition of quark deconfinement in accretion-induced collapse of white dwarfs

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We present the first general relativistic, neutrino-radiation hydrodynamics simulations of accretion-induced collapse (AIC) extending to seconds after core bounce, using realistic hadron–quark hybrid equations of state (EOSs). A first-order QCD phase transition (PT) triggers a second dynamical collapse and the formation of a quasistable protohybrid star (PHS) with a deconfined quark core and a distinctive second neutrino burst. We find that the thermally suppressed onset of the mixed phase allows low-mass protoneutron stars to enter the hadron-quark mixed phase during long-term evolution, even for hybrid EOSs with high onset densities. In contrast to core-collapse supernovae (CCSNe), AIC models exhibit a tightly constrained onset mass with minimal EOS dependence, owing to the absence of a massive envelope and thus the reduced postbounce accretion. This enhances the sensitivity of neutrino observables in AIC to hybrid EOS properties. We establish empirical relations between PT onset density and neutrino signatures, revealing a distinct behavior in AIC not seen in CCSNe. Our results suggest that a single Galactic AIC neutrino detection could place strong constraints on QCD PT thresholds, hybrid EOS characteristics, and the existence of PHSs. PT in AIC may also produce gravitational waves, gamma-ray bursts, and *r*-process elements, motivating multidimensional simulations with rotation, magnetic fields, and improved microphysics for realistic multimessenger predictions.

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I. INTRODUCTION

White dwarfs (WDs) are compact stellar remnants supported by electron degeneracy pressure. The outcome of an accreting WD is determined by its chemical composition, initial mass, and accretion rate. Carbon-oxygen WDs with masses of $\lesssim 1.2M_{\odot}$ and high accretion rates are prone to carbon deflagration, leading to a type Ia supernova that leaves no remnant [1–4]. In contrast, oxygen-neon WDs with masses of $\gtrsim 1.2M_{\odot}$ and slower accretion rates can continue to gain mass until they reach the Chandrasekhar limit beyond $1.44M_{\odot}$ [5]. This triggers a gravitational collapse due to electron capture, followed by a core bounce, resulting in the formation of a protoneutron star (PNS) remnant—a phenomenon known as accretion-induced

collapse (AIC) and the resulting explosion, known as an electron-capture supernova (SN) [6,7].

AIC can occur via two main channels: mass transfer from a companion star to a WD in a binary system [8–10], or binary WD merger [11–16]. The resulting super-Chandrasekhar WDs are often rapidly rotating [17–19] and highly magnetized [20–24] in both channels. The collapse of such progenitors is expected to produce multimessenger signals, including gravitational waves and neutrinos [25–28], and to explain the origin of millisecond pulsars [9]. The ejecta mass driven by neutrino heating can potentially power rapid neutron capture (*r*-process) nucleosynthesis [29–33]. Furthermore, these collapses can potentially power gamma-ray bursts [32,34–36], fast radio bursts [37,38] and fast blue optical transients [39,40].

The galactic rate of AIC events is estimated to be $(0.3 - 0.9) \times 10^{-3} \text{ yr}^{-1}$ from mass transfer onto a WD [10], and $(1.4 - 8.9) \times 10^{-3} \text{ yr}^{-1}$ from binary WD mergers [14,15]. Although the occurrence rates remains uncertain, there is observational evidence of AIC [41–43]. Moreover, several systems have been proposed as potential AIC progenitors or remnants, based on their total mass, magnetic field strength, or binary configuration [44–50].

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The progenitor mass of AIC depends on the formation channel. AIC resulting from mass transfer onto a single WD typically occurs near the Chandrasekhar limit [10], whereas AIC from binary WD mergers can reach up to $\sim 2.5M_{\odot}$ [14,15]. This is substantially smaller than the progenitor mass of SN systems with a massive envelope, such as core-collapse supernovae (CCSNe), which range from $8 - 100M_{\odot}$ [51]. Besides, the low-mass nature leads to a considerable astrophysical interest owing to its characteristically low-mass progenitors. It is expected to explain the formation of low-mass neutron stars [52], and low-mass millisecond pulsars or x-ray binaries [53]. The ejecta mass, which depends on the magnetic field strength and rotation period [32,54], is bounded by the difference between the initial WD and the final PNS masses, typically ranging up to $\sim 10^{-1}M_{\odot}$ [28,32,54]. This is generally smaller than that of CCSNe associated with the collapse of massive stars, due to the low-mass progenitors of AIC. Moreover, the absence of a massive envelope allows the ejecta to maintain their initial high velocities.

The absence of a massive envelope and the low-mass nature of the progenitor of an AIC system lead to a short accretion phase after core bounce. Thereafter, neutrino emission from the surface carries away thermal energy, and the resulting loss of thermal pressure drives the contraction of the PNS [55,56]. This process, known as Kelvin-Helmholtz cooling, proceeds over a timescale of several seconds. Consequently, the density and temperature at the core of the PNS can increase continuously, though at a slower rate than during the accretion phase. Several-second-scale simulations of the PNS cooling phase have shown that the central temperature can rise to 40–60 MeV within the first ten seconds, before eventually decreasing [55,57–59]. This long-term evolution naturally raises the question of whether the conditions within AIC systems become suitable for triggering a quantum chromodynamics (QCD) phase transition (PT).

A conservative estimate is to utilize the phase diagram of a hadron-quark hybrid equation of state (EOS) to examine whether the onset of the hadron-quark mixed phase can be reached by a typical PNS near nuclear saturation density ρ_{sat} and temperatures of ~ 10 – 20 MeV. For some hybrid EOSs with onset conditions near ρ_{sat} [60–62], it is reasonable to expect that QCD PT can take place in the low-mass AIC system. A key remaining question is whether the thermodynamic conditions in AIC systems can reach the PT onset threshold for certain hybrid EOSs, particularly in cases where the onset density of PT is suppressed at higher temperatures [61–66].

In similar compact stellar systems such as CCSNe, hybrid EOSs featuring a first-order QCD PT have been shown to induce a second collapse, halted by the formation of a stiffer deconfined quark core, and to generate a secondary hydrodynamical shock accompanied by a neutrino burst [67–76]. Moreover, such a collapse can produce a strong burst of

gravitational waves [69], and the explosion driven by the QCD PT has been proposed as a potential r-process site [77], with its properties closely linked to the characteristics of the hybrid EOS [76,78]. These findings highlight the astrophysical and observational significance of studying a first order QCD PT in the AIC system. The properties of the PT, such as the onset conditions and the location of the critical point, are a central topic in nuclear astrophysics for understanding the behavior of dense matter. These aspects have been extensively investigated through theoretical studies of dense matter EOS [79–81], gravitational-wave astronomy [82–84], and observation [85,86]. Since the primary difference between AIC and CCSNe is the absence of a massive envelope, it is essential to investigate how the dynamics of a QCD PT may unfold in the AIC scenario, where the lack of an extended envelope could potentially alter the resulting evolution and observables.

In this work, we present the first seconds-long general relativistic neutrino-radiation simulations of the AIC systems that employ realistic hadron-quark hybrid EOSs to explore the influence first-order QCD PT, the existence of deconfined quark matter and the properties of hybrid stars.

Unlike the isolated neutron star setups adopted in Refs. [87,88], our super-Chandrasekhar mass WD model follows the collapse and formation of a PNS embedded in an accreting envelope. This leads to a strongly nonequilibrium environment with evolving density, temperature, and electron fraction regulated by neutrino transport, which is absent in equilibrium neutron star configurations. As a result, the PT in our work is determined self-consistently by the collapse dynamics rather than imposed on a preexisting neutron star.

This paper is structured as follows. In Sec. II, we present our numerical setup. Our main results are presented in Sec. III, followed by a discussion of future research directions in Sec. IV.

II. NUMERICAL METHODS AND SETUPS

We use a set of hadron-quark hybrid EOSs, known as relativistic density functional (RDF) EOSs, where both quark and hadron matter are described within the RDF formalism [61,89]. In the RDF formalism, both the hadronic and quark models are described by a set of coupling parameters that determine the underlying interaction potentials [61]. A first-order PT connects the hadronic phase to the stiffer deconfined quark phase, forming mixed nuclear-quark matter constructed via Maxwell reconstruction under the assumption of mechanical and chemical equilibrium. For each set of coupling parameters, the latent heat is set by the pressure-density jump between hadronic and quark branches, and controlled by quark couplings tuned to fix the onset density and density jump [61]. The RDF EOSs incorporate varying onset densities and latent heats, with deconfined quark matter modeled using a string-flip approach. These EOSs are consistent with astrophysical

TABLE I. The hybrid EOSs employed in our simulations. ρ_{onset} and M_{onset} denote the onset density and onset gravitational mass of the PT starting to reach the mixed phase, where these three quantities are defined under $T = 0$ MeV and beta equilibrium [76]). The maximum gravitational mass of the nonrotating model is denoted as M_{max} for each EOS.

EOS	$\rho_{\text{onset}} [10^{14} \text{ g cm}^{-3}]$	$M_{\text{max}} [M_{\odot}]$	$M_{\text{onset}} [M_{\odot}]$
RDF-1.9	4.6	2.16	0.81
RDF-1.2	7.3	2.15	1.37
RDF-1.5	8.2	2.03	1.46
RDF-1.1	8.8	2.13	1.55

constraints, including precise high-mass neutron star measurements [90] and recent NICER results [91]. These EOSs have been recently utilized in astrophysical simulations [65,71–73,76]. Detailed properties are listed in Table I and can also be found in Refs. [61,76].

Since a given set of coupling parameters does not produce systematic changes in key observables, such as the onset density and density jump (latent heat) of the hadron-quark PT, we select four hybrid EOSs with notably different onset densities to simplify the comparison. The corresponding latent heats do not vary in an equally systematic manner, so we restrict our systematic analysis to the onset conditions.

We focus on the double degenerate binary merger channel for AIC [11–16], in which the merger of two WDs—including double ONeMg, ONeMg-CO, or double CO configurations—produces a super-Chandrasekhar remnant that undergoes AIC rather than thermonuclear explosion. In particular, for double CO WD mergers, nonstandard evolutionary pathways involving off center carbon ignition can transform the merger remnant into an ONe-core configuration that may subsequently collapse to a neutron star [15,16].

All progenitors are generated using RNS [92]. They are all spherically symmetric, nonrotating, and in beta equilibrium. The central energy density ϵ_c/c^2 is 10^9 g cm^{-3} . The gravitational mass and baryonic mass are $1.53M_{\odot}$ and $1.54M_{\odot}$, respectively, with a radius of 2800 km. The use of massive WD initial models is commonly adopted in AIC studies [28,30–32,93]. Although rotation is known to significantly affect the dynamics of AIC [28,31,32], we neglect rotational degrees of freedom in this study to isolate the effects of the hadron-quark EOS on collapse dynamics and neutrino signals. The impact of rotation will be explored in future work. Since RNS does not solve for temperature, all progenitors generated are cold and require additional construction before the evolution.

Our simulations employ the general relativistic neutrino-radiation hydrodynamics code Gmunu [94–98], which solves Einstein’s field equations under the conformal flatness condition [94,99,100].

The initial progenitor is constructed with a fixed temperature of 0.01 MeV corresponding to the minimum of the EOS table and an electron fraction of $Y_e = 0.5$ in Gmunu. To model the temperature profile of a hot white dwarf prior to collapse, we superimpose a temperature profile using the relation $T = T_c(\rho/\rho_c)^{0.35}$ with $T_c = 5 \times 10^9$ K corresponding to ~ 0.43 MeV following the approach of Refs. [30,32,93]. The addition of this temperature profile together with the specific metric initialization reduces the gravitational mass to $M_g = 1.52M_{\odot}$. This inconsistency in M_g arises because Gmunu adopts the metric initialization method of Ref. [99], which fixes the initial primitive variables after the thermal profile is imposed. However the metric initialization should instead be performed by fixing the conformal factor and the conserved variables as done for the hot hypermassive neutron star postmerger simulation in Ref. [100], which will be adopted in the future work.

We adopt the energy-integrated two-moment neutrino transport scheme [101,102], and the implicit-explicit Runge-Kutta method IMEX-SSP2(2,2,2) [103] for time integration. Details of the implementation and comparisons with similar numerical schemes [104,105] are provided in the appendix of [102]. For the neutrino microphysics, we employ the neutrino library Weakhub [98] and consider the conventional set of weak interactions defined in [98] but excluding inelastic neutrino-electron scattering. Neutrino opacities are tabulated after averaging in energy space [102], where the neutrino energy is logarithmically discretized into 20 bins within the range [0.1, 400] MeV. The emissivities are recalculated on the fly using Kirchhoff’s law.

The collapse of WD is triggered by the electron capture, in which the electron capture rates on heavy nuclei are uncertain [106] during the collapse of WD. Our initial models are constructed with a central density of $\rho_c = 10^9 \text{ g cm}^{-3}$. Also, energy-integrated neutrino transport schemes neglect energy-coupled interactions, such as neutrino-electron inelastic scattering, which are crucial during the collapse phase. To bridge this modeling gap while maintaining computational efficiency, we apply an effective deleptonization scheme [107] to mimic the proton fraction Y_p profile as a function of rest-mass density ρ during the collapsing phase, which is widely used in AIC simulations [27,28,32]. We then switch to the two-moment neutrino transport when the core bounces (defined as when the matter specific entropy $s \geq 3 k_b/\text{baryon}$ in the core region), and enable the coupling between neutrinos and the fluid before core bounce.

The computational domain extends from 0 to 3×10^3 km at the outer boundary, with a radial resolution of $N_r = 128$. We allow a 10 adaptive mesh refinement (AMR) level and the finest grid spacing at the center of the star is 45.8 m. To demonstrate the robustness of our results, we present a comparison with a recent CCSN PT simulation [69] and

TABLE II. Various quantities extracted at the moment of first reaching the mixed nuclear-quark phase (columns in the left part) and 300 ms after the first entry into the deconfined quark phase (columns in the right part) in our AIC simulations, using different RDF EOS. t_{mixed} , $\rho_{\text{mixed,c}}$, $T_{\text{mixed,c}}$, M_{mixed}^b and $\mathcal{L}_{\text{mixed,tot}}$ denote the time after the first core bounce at $t = t_b$, central density, central temperature, the enclosed baryon mass of the PNS and the total neutrino luminosity for the system first reaching the mixed nuclear-quark phase. t_Q and $M_{t_Q}^b$ are the time with respect to t_b and the enclosed baryon mass of the PHS when the PHS first reaching the deconfined quark phase. Baryon mass (radius) of the quark core, the mixed-phase mantle (MM) and the new-born protohybrid star (PHS), are expressed as M_Q^b (r_Q), M_{MM}^b (r_{MM}) and M_{PHS}^b (r_{PHS}). The values of neutrino luminosity and averaged energy are evaluated at $r = 500$ km. The values of the second neutrino burst are represented as $\mathcal{L}_{\nu_e}$ and $\langle \epsilon_{\nu_e} \rangle$, respectively for electron neutrino ν_e , electron antineutrino $\bar{\nu}_e$ and a collective species describing heavy lepton neutrinos and their antineutrinos as ν_x . Note that the PNS or PHS are defined as the regions of $\rho > 10^{11} \text{ g cm}^{-3}$.

EOS	t_{mixed} [s]	$\rho_{\text{mixed,c}}$ [$10^{14} \text{ g cm}^{-3}$]	$T_{\text{mixed,c}}$ [MeV]	M_{mixed}^b [$M_\odot$]	$\mathcal{L}_{\text{mixed,tot}}$ [$10^{51} \text{ erg s}^{-1}$]	t_Q [s]	$M_{t_Q}^b$ [$M_\odot$]	M_Q^b [$M_\odot$]	M_{MM}^b [$M_\odot$]	M_{PHS}^b [$M_\odot$]	r_Q [km]	r_{MM} [km]	r_{PHS} [km]	$\mathcal{L}_{\nu_e}$ [$10^{51} \text{ erg s}^{-1}$]	$\mathcal{L}_{\bar{\nu}_e}$ [$10^{51} \text{ erg s}^{-1}$]	$\mathcal{L}_{\nu_x}$ [$10^{51} \text{ erg s}^{-1}$]	$\langle \epsilon_{\nu_e} \rangle$ [MeV]	$\langle \epsilon_{\bar{\nu}_e} \rangle$ [MeV]	$\langle \epsilon_{\nu_x} \rangle$ [MeV]
RDF-1.9	0.58	4.1	17.8	1.53364	41.9	1.11	1.53443	1.082	0.15	1.524	5.93	6.66	12.7	148	837	501	28.8	59.6	104.7
RDF-1.2	1.56	5.3	23.4	1.53459	22.6	2.91	1.53397	0.92	0.23	1.531	5.42	6.34	10.7	194	412	187	44	50.5	91.8
RDF-1.5	3.06	5.94	26.6	1.53468	14.4	4.27	1.53346	0.875	0.19	1.533	4.92	5.61	9.64	49.8	65.2	21.2	31.8	34.1	46.9
RDF-1.1	3.32	6.03	27.5	1.53469	14.2	6.44	1.53397	0.288	0.61	1.533	4.14	5.29	9.82	6.59	16	8.03	20.8	24.3	24.2

include a resolution dependence test of our AIC setup in Appendix A and B respectively.

III. RESULTS

The unstable initial WD undergoes a collapse and an electron capture following the effective deleptonization scheme, which reduces the electron degeneracy pressure in the WD. Electron capture becomes significant at $\sim 5 \times 10^9 \text{ g cm}^{-3}$ and triggers a gravitational collapse with a rapid increase in central density ρ_c . When ρ_c exceed nuclear saturation density ($\rho_c \gtrsim 10^{14} \text{ g cm}^{-3}$), hadronic matter stiffens due to the strong nuclear force and halt the collapse, lead to a core bounce and an outward shock at $t = t_b$. This produces a ν_e burst with $\mathcal{L}_{\nu_e} \sim 10^{53} \text{ erg s}^{-1}$ over ~ 10 ms, which is qualitatively consistent with the result in Refs. [28,31,32] and detailed in Appendix C.

The PNS gradually loses energy through weaker neutrino emissions across all species, maintaining $\mathcal{L}_{\nu_e} \sim 10^{51} \text{ erg s}^{-1}$ for several seconds. This sustained emission depletes thermal energy and pressure, leading to gradual contraction via Kelvin-Helmholtz cooling, which slowly increases ρ_c and T_c .

The contraction accelerates upon reaching the onset of a first-order PT, forming mixed hadron-quark phase matter, i.e. quark fraction $Y_q > 0$, which softens the EOS. The left columns of Table II show the time t_{mixed} , central density $\rho_{\text{mixed,c}}$, central temperature $T_{\text{mixed,c}}$, the enclosed baryonic mass of the PNS M_{mixed}^b and the total neutrino luminosity $\mathcal{L}_{\text{mixed,tot}}$ when the system first reaches the mixed phase with $Y_q > 0$.

For all models, M_{mixed}^b differs by under 0.07% across models, indicating that the accretion phase is effectively ended and a steady PNS has formed at this point (see discussion in Fig. 2). The small variation in M_{mixed}^b

indicates a degree of universality in the PT onset due to the absence of a massive envelope.

The negligible variation in onset mass distinguishes PT in AIC from that in CCSNe, where the presence of a massive envelope leads to continuous mass accretion and possibly different onset mass condition. These considerations motivate an investigation of the potential difference in PT dynamics; we perform a set of CCSN simulations using the lowest-mass progenitor model, *s12* from Ref. [51], with the same set of hybrid EOSs. We focus the discussion of PT in AIC in the main text, and thus the numerical setup and detailed comparison are provided in Appendix D.

Due to high T_c and low Y_p at the PNS core, the modified onset density $\rho_{\text{mixed,c}}$ is lower than ρ_{onset} in Table I, which is defined under $T = 0$ MeV and beta equilibrium [76]. For example, ρ_{mixed} is 30% (10%) lower than ρ_{onset} in model RDF-1.1 (RDF-1.9), corresponding to $T_{\text{mixed,c}}$ of 27.5 MeV (17.8 MeV) at $t = t_{\text{mixed}}$.

As contraction accelerates at the onset of a first-order PT, the emergence of a mixed phase with nonzero quark fraction ($Y_q > 0$) further softens the EOS. In our case, this PT-induced contraction proceeds concurrently with the Kelvin-Helmholtz cooling, where models with higher onset densities receive less assistance from cooling, as indicated by the substantial decrease in $\mathcal{L}_{\text{mixed,tot}}$.

Figure 1 illustrates the evolution of central density ρ_c (solid) and quark fraction $Y_{q,c}$ (dotted) through three phases: the mixed phase ($Y_{q,c} < 1$), the onset of deconfined quark matter ($Y_{q,c} = 1$) at t_Q (first entry into the deconfined quark phase relative to t_b , see Table II), and the quasistable PHS phase.

In the first phase, increasing Y_q accelerates collapse as the EOS softens. A greater fraction of mixed-phase matter leads to faster collapse, meaning higher $\rho_{\text{mixed,c}}$ results in a

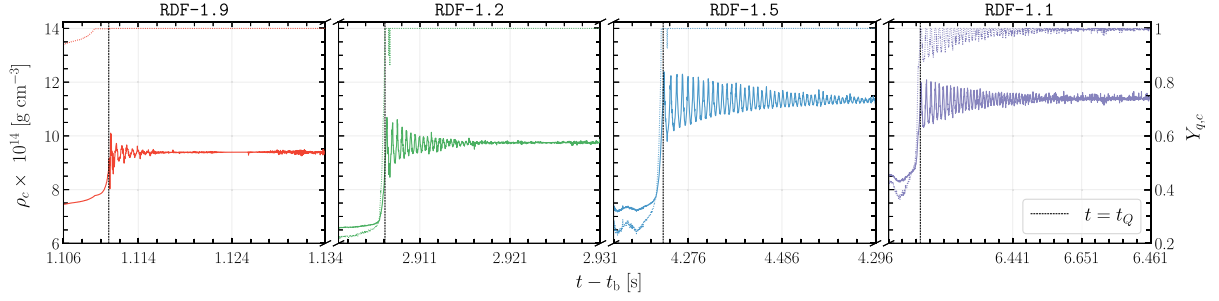


FIG. 1. Time evolution of the central density ρ_c (solid lines) and central quark fraction $Y_{q,c}$ (dotted lines) after the first core bounce for all models. The second bounce for $t - t_Q < 23$ ms are shown. Dashed lines correspond to the time at $t = t_Q$.

longer time to reach t_Q . As the core approaches the stiffer deconfined quark phase at $t \approx t_Q$, a second dynamical collapse happens within 0.5 ms. Notably, the PHS collapse with enclosed baryon mass at t_Q with $M_{t_Q}^b \sim 1.53M_\odot$ across all models (See Table II), highlighting that PT-induced contraction in AIC is not subject to mass accretion, enhancing its sensitivity to properties of the hybrid EOS, in contrast to the behavior of mixed phase observed in CCSNe (see Appendix D).

For all RDF EOSs, the system's small baryonic mass ($\sim 1.54M_\odot$) and the stiffness of deconfined quark matter disfavor black hole formation. Instead, the stiff quark core halts the infall of mixed-phase material, generating a second outward-propagating shock. This shock, at high density and temperature, converts infalling hadronic matter into the mixed phase and accelerates as it reaches the lower-density hadronic envelope near the PNS surface.

The second shock energizes the hadronic envelope, enhances e^+ captures, and produces a second neutrino burst (see detailed discussion on the second burst later). This second burst from PT is consistent with Appendixes A and D and with the previous studies of PT in CCSNe [67,69–73,76].

A quasistable PHS forms in each case, comprising a quark core, a mixed-phase mantle, and a residual hadronic envelope with baryonic masses M_Q^b , M_{mixed}^b , and M_{PHS}^b , respectively, as listed in Table II (right columns). The overcompression of the quark core before the second bounce, combined with dynamical changes in quark and mixed-phase matter, induces strong perturbations in the newborn PHS, leading to rapid oscillations that damp out within ~ 23 ms in Fig. 1.

To examine the dynamics around the second bounce, we focus on the RDF-1.5 EOS case. The upper panel of Fig. 2 shows the specific entropy s evolution from $t - t_b$ to 300 ms after t_Q . The first core bounce generates a shock that expands to ~ 90 km within 5 ms. By $t - t_b \sim 10$ –100 ms, the shock loses energy, stalls at ~ 100 km, and forms an accreting envelope with $s \sim 7$ –15 k_b /baryon.

After core bounce, the PNS mass increases from $0.92M_\odot$ at $t = t_b$ to $1.534M_\odot$ by $t - t_b \sim 600$ ms due to the

accretion of infalling shocked matter and material from the outer envelope.

As accretion diminishes, the system reaches a steady state with a baryonic mass of $1.535M_\odot$ at $t - t_b \approx 1.2$ s. Despite entering the mixed phase, the PNS continues to contract (see yellow line in Fig. 2), reinforcing that neutrino cooling and mixed-phase matter soften the EOS. Additionally, the high-entropy matter ($s \sim 25$ –50 k_b /baryon) observed at 30–100 km after $t - t_b \sim 1$ s arises from interactions between the accretion envelope and fallback of ejected matter from the first shock.

Immediately after $t = t_Q$, the newly formed PHS core undergoes a second bounce, generating a shock that propagates to ~ 500 km in 2.4 ms at $\sim 0.5c$ (see Fig. 2, lower panel). As it pushes against dense matter beyond the accretion envelope, heating it to ~ 50 k_b /baryon, it loses

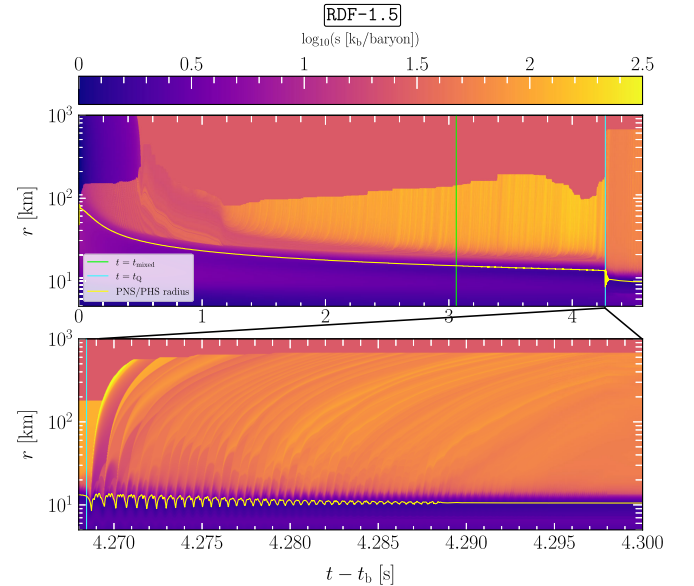


FIG. 2. Specific entropy s profile of the model with RDF-1.5 EOS. Top panel: from $t = t_b$ to 300 ms after t_Q . Bottom panel: from $t = t_Q$ to 32 ms after t_Q . The vertical lines indicate time moments of $t = t_{\text{mixed}}$ (lime) and $t = t_Q$ (cyan). The yellow solid represents the radius of PNS or PHS.

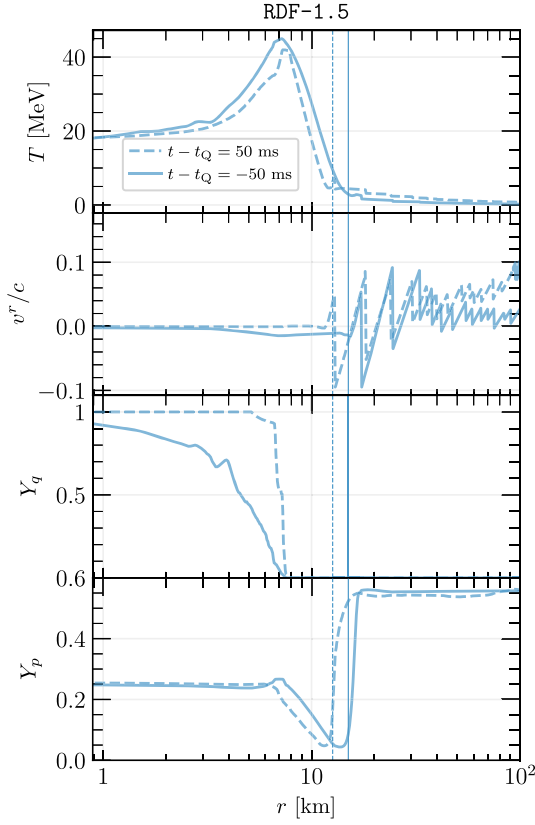


FIG. 3. Radial snapshots of temperature T , radial velocity v^r/c , proton fraction Y_p , and quark fraction Y_q of RDF-1.5 model at 50 ms before (solid) and 50 ms after (dashed) the time moment of t_Q . Vertical lines correspond to the density at $\rho = 10^{11} \text{ g cm}^{-3}$.

energy and stalls at $\sim 500 \text{ km}$. Before stalling, PHS oscillations produce multiple weaker shocks at $\sim 0.25c$ that continue pushing matter outward, expanding the envelope to $\sim 600 \text{ km}$ within 4 ms. Since the accreting envelope in AIC systems is only $\sim 0.1\%$ of the total mass, the second shock loses less energy and propagates beyond the first shock stall ($\sim 100 \text{ km}$), reaching at least 500 km (RDF-1.1) and up to 1500 km (RDF-1.9). Following the second bounce, the PHS rapidly contracts, with the radius shrinking by $\sim 20\%$ at 50 ms after t_Q and continuing to contract gradually. This shock, propagating through the low-proton-fraction region (see the bottom panel of Fig. 3), generates a second neutrino burst.

Figure 3 shows the radial hydrodynamical profiles at 50 ms before and after t_Q for the RDF-1.5 model. The vertical lines indicate the radii of PNS (dashed) and PHS (solid), respectively. Note that by $t - t_Q = 50 \text{ ms}$, the PHS attains a quasisteady state and is much more compact (see also Fig. 1).

The PHS exhibits a cooler core compared to the PNS at $t - t_Q = -50 \text{ ms}$, as the thermal energy is converted to latent heat during the formation of deconfined quark matter.

In contrast, the shocked envelope outside the PHS (for $r > 20 \text{ km}$) is slightly hotter.

Unlike the PNS at $t - t_Q = -50 \text{ ms}$, which features oscillating matter in its envelope, the quasisteady PHS continues to eject materials with $v^r > 0$ as long as 300 ms. In these outer layers, the velocity can exceed $0.1c$ at $r > 100 \text{ km}$. This outflow is driven by the heavily perturbed PHS, where neutrino interactions become active again in the shocked regions. These neutrinos deposit energy via reabsorption in the outer layers, leading to the ejection of about 0.1% of the PHS mass by $t - t_Q = 300 \text{ ms}$.

The fractions of deconfined quark matter, mixed-phase matter and hadronic matter in the PHS are highly nonlinear and depend on the EOS properties that determine the collapse dynamics and the strength of the second bounce (the properties of PHS can be found in the right columns of Table II). In all models, the mixed-phase mantle is sandwiched between the hadronic envelope and deconfined-quark core with a small size of at most 2 km due to its softness among all cases. Even a special case of RDF-1.1, where the mass of the mixed-phase mantle exceeds that of the quark core, follows this pattern.

Particular attention should be paid to this special case of RDF-1.1. Its PNS collapse driven by the PT is weaker and slower due to its high $\rho_{\text{mixed},c}$ and low $\mathcal{L}_{\text{mixed},\text{tot}}$, leading to a longer timescale of cooling and thus slower velocity of collapsing matter before t_Q . The slower contraction is reflected in a small but notable decrease in both $Y_{q,c}$ and ρ_c at $t - t_Q \approx -3.8 \text{ ms}$ in Fig. 1. These result in a much lighter quark core surrounded by a significantly more mixed-phase mantle. The result of RDF-1.1 case is nearly a critical case where the progenitor mass with $1.51M_\odot$ (gravitational mass) is close to the onset mass of the EOS with $1.57M_\odot$ shown in Table I, and it is also the reason for oscillating in $Y_{q,c}$ after the second bounce.

Two additional robust features appear in all models. First, a cavity with a low proton fraction ($Y_p \sim 0.1$) forms at about $t - t_b \sim 10 \text{ ms}$, coinciding with the ν_e burst shortly after the core bounce, and this cavity persists until the PHS forms (see the Y_p profile in Fig. 3). The second shock, originating from the deconfined quark core beneath the cavity, passes through this low Y_p region, emitting primarily $\bar{\nu}_e$ via positron (e^+) capture. Consequently, the second burst is predominantly composed of $\bar{\nu}_e$ for all EOSs, similar to the behavior reported in CCSN simulations with a QCD PT [68,76]. In contrast, the first neutrino burst, which results from the first bounce, is dominated by ν_e emission due to the proton-rich environment of the newly formed PNS.

Finally, the high temperature of the PHS leads to a greater integrated emission of ν_x compared to ν_e during the second burst. Furthermore, the second burst generally produces neutrinos with higher average energy across all species compared to the first burst.

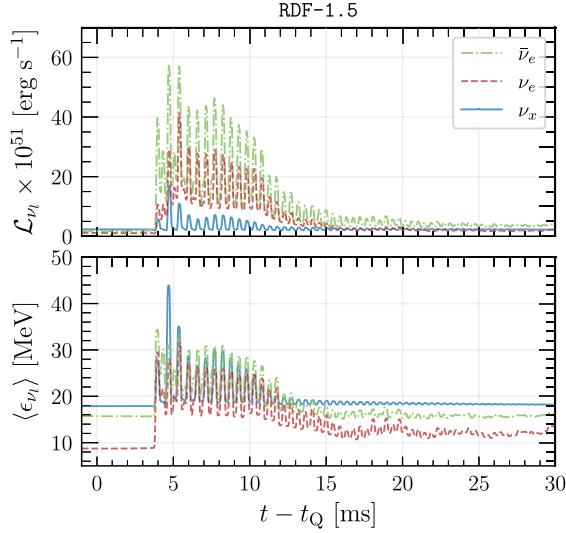


FIG. 4. Time evolution of neutrino luminosity $\mathcal{L}_{\nu_i}$ and average energy $\langle \epsilon_{\nu_i} \rangle$ of the second neutrino burst formed in the case with RDF-1.5 EOS for ν_e (red dashed), $\bar{\nu}_e$ (green dash-dotted) and ν_x (blue solid), respectively.

Figure 4 presents the neutrino luminosity, $\mathcal{L}_{\nu_i}$, and average energy, $\langle \epsilon_{\nu_i} \rangle$, for all species during the second millisecond neutrino burst in the RDF-1.5 model. The burst, with luminosities on the order of $10^{52} \text{ erg s}^{-1}$, is triggered by the second shock as it propagates through the PHS neutrinospheres and is observed around 4 ms after t_Q .

At the onset of the burst, neutrino luminosities and average energies for all species reach a peak, as summarized in Table II (right columns). The burst is dominated by $\bar{\nu}_e$ due to positron (e^+) capture in the low-proton-fraction cavity during shock propagation. Also, the stronger second shock produces neutrinos with higher average energy across all species compared to the first shock. A comparison of the first and second bursts is detailed in Appendix C.

These quantities then oscillate for ~ 30 ms, corresponding to the oscillation and damping of the PHS following the second bounce (Fig. 1). This oscillatory behavior in the second burst, lasting 5–30 ms, is consistently observed across all models, and is also indicated in the CCSNe simulations in previous studies [70,108].

The second neutrino burst not only confirms a strong first-order PT and PHS formation but also provides a key observational probe. The time interval between the two bursts, approximately t_Q , can potentially constrain the onset density and PT properties of the EOS. As shown in Table II, there are small variations in $\rho_{\text{mixed,c}}$ across different EOSs.

Figure 5 presents empirical relations between t_Q (top panel), total neutrino luminosity $\mathcal{L}_{\text{tot}}$ (middle panel), and $\langle \epsilon_{\bar{\nu}_e} \rangle$ (bottom panel) as functions of $\rho_{\text{mixed,c}}$. Results from AIC and CCSN simulations in Appendix D are shown for comparison to highlight systematic differences due to the presence of a massive envelope. The top panel shows that

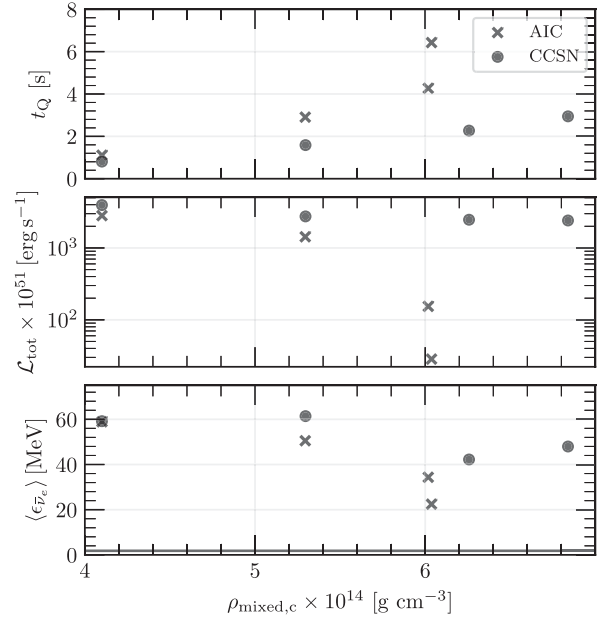


FIG. 5. Empirical relations of t_Q , the total neutrino luminosity $\mathcal{L}_{\text{tot}}$ and $\langle \epsilon_{\bar{\nu}_e} \rangle$ as functions of $\rho_{\text{mixed,c}}$. Relations from PT in CCSN and AIC are marked with circles and crosses respectively. The blue horizontal line in the bottom panel indicates 1.8 MeV inverse beta decay threshold in water Cherenkov detectors [109–112].

t_Q increases significantly with $\rho_{\text{mixed,c}}$, unlike QCD PT simulations in CCSNe.

The discrepancy arises because the accreting and infalling mass is negligible compared to that in CCSNe in AIC systems. As a result, models with higher PT onset densities require stronger neutrino cooling to trigger the PT-induced contraction, leading to a significantly longer t_Q , as discussed in detail above.

For these reasons, different relationships between neutrino observables and $\rho_{\text{mixed,c}}$ emerge (see the middle and bottom panels of Fig. 5).

First, both t_Q and $\langle \epsilon_{\bar{\nu}_e} \rangle$ exhibit greater sensitivity to variations in $\rho_{\text{mixed,c}}$ in AIC systems, resulting in a wider dynamic range of response. In particular, $\langle \epsilon_{\bar{\nu}_e} \rangle$ in AIC shows a clearer and more rapid decrease compared to CCSN. Despite the substantial drop in $\langle \epsilon_{\bar{\nu}_e} \rangle$ for AIC, it remains above the 1.8 MeV inverse beta decay threshold in water Cherenkov detectors [109–112], showing that neutrino bursts remain potentially detectable for all AIC models.

Moreover, the relation between $\mathcal{L}_{\text{tot}}$ and $\rho_{\text{mixed,c}}$ unambiguously shows a transition from linear dependence in CCSN to exponential dependence in the AIC system. Although this dependence implies a rapid decline in detectability with increasing $\rho_{\text{mixed,c}}$, it establishes a distinct and potentially powerful relation for constraining the onset density of the first-order PT. Although a $\sim 10^{51} \text{ erg s}^{-1}$ neutrino burst from an AIC event occurring within our

Galaxy may remain detectable by supernova-dedicated neutrino observatories, such as Super-Kamiokande and Hyper-Kamiokande [111,112], a comprehensive assessment of detectability across current and next-generation detectors will be conducted in future work [113].

These empirical relations directly link potentially detectable neutrino signals to $\rho_{\text{mixed,c}}$ in AIC systems, offering distinct advantages for constraining the onset density of the first-order PT. Notably, these relations demonstrate that the absence of a massive envelope in AIC substantially enhances its sensitivity to the properties of the hybrid EOS compared to CCSN.

Due to the advantageous second-long time separation between the two neutrino bursts in AIC systems, along with their less massive accretion envelopes and a narrower total mass range of $1.44 - 2.5M_{\odot}$ [14,15], in contrast to the broader $8 - 100M_{\odot}$ range in CCSNe [51], exploring first-order PT in AIC shows a stronger dependence on the onset properties of hybrid EOS, particularly in neutrino observables, than PT in CCSNe. The strong EOS sensitivity in AIC could enable exceptionally tight constraints on PT and QCD matter properties from the detection of a single Galactic AIC event.

IV. DISCUSSIONS AND CONCLUSIONS

We present the first seconds-long general relativistic neutrino-radiation simulations of AIC with realistic hadron-quark hybrid EOSs, evolving the collapse of WDs. A first-order PT triggers a second dynamical collapse and the formation of a quasistable PHS with a deconfined quark core and a distinct second neutrino burst.

Following core bounce, the PNS undergoes slow contraction mainly due to neutrino cooling, with the appearance of mixed nuclear-quark matter further softening the core and slightly accelerating collapse. We find that the thermally suppressed onset density of mixed phase allows a low-mass PNS to enter the mixed phase during a long-term and second-scale evolution for hybrid EOS characterized by a high onset density for the mixed phase. We further observe that the absence of a massive envelope allows all AIC models to exhibit a tightly constrained onset mass in both the mixed nuclear-quark phase and deconfined quark phase, showing minimal hybrid EOS dependence. The absence of sustained accretion indicates that, compared to PT in CCSNe, AIC receives substantially less support from mass accretion during the mixed phase. As a consequence, the time spent in the mixed phase becomes more strongly dependent on the properties of the hadron-quark hybrid EOS. Therefore, the low-mass nature of AIC progenitors enhances its sensitivity to the properties of the hybrid EOS compared to CCSNe.

The transition to a stiffer deconfined quark core halts the collapse, generating a second shock breakout and a $\bar{\nu}_e$ -dominated burst with higher average neutrino energy across all species compared to the first burst. In all cases, a

stable PHS forms, consisting of a deconfined quark core, a mixed-phase mantle, and an outer hadronic envelope. Although the formation of hybrid quark stars near the Chandrasekhar mass is possible for hybrid EOSs with low onset densities [61], such scenarios may conflict with observations of pulsars with comparable masses [114,115]. Our results indicate that if the cooled down hybrid star structure persists, certain PT models could be ruled out.

AIC provides a promising avenue for constraining PT thresholds and hybrid EOS properties. Unlike CCSNe, its narrower mass range and a less massive accreting envelope makes them ideal for studying QCD PT and PHS properties through observational signatures. We establish an empirical relation between PT onset density and neutrino observables, particularly the time delay between the two bursts, total neutrino luminosity and the average energy as a function of the PT onset density. This allows a direct comparison with the corresponding relation in CCSN and reveals a new, distinct relation in AIC, primarily driven by the low-mass nature of its progenitor.

Although the rate of AIC ($\lesssim 10^{-3} \text{ yr}^{-1}$) is lower than that of CCSN ($\sim 10^{-2} \text{ yr}^{-1}$), our results suggest that a single Galactic AIC neutrino detection could impose more precise constraints on PT thresholds, hybrid EOS properties, and the existence of hybrid stars. These energetic outbursts are also a promising new source of GW emission, gamma-ray bursts, and r -process nucleosynthesis [77].

It is observed that PT in CCSN, and the associated neutrino observables, are sensitive to the different physics and numerical setup (see Appendix A). Therefore, future work of PT in AIC should incorporate multidimensional effects, including accretion instabilities [31]. Rotation also plays a crucial role in AIC [28,32], where it can potentially regulate PNS cooling and contraction, modify the second collapse dynamics, and thereby shape the detailed evolution of the PT. In multidimensional simulations convection and anisotropic temperature profiles can also play an essential role in the formation and transport of mixed hadron quark matter as observed in multidimensional simulations of neutron star mergers [116,117]. More advanced microphysics will affect the Kelvin-Helmholtz cooling strength and, hence, alter the neutrino observables such as t_Q [118–120]. The strength and geometry of magnetic fields are known to affect AIC dynamics [32,121], and their potential influence on the second collapse will be crucial. A comprehensive study of the detectability of the neutrino signal in the current and next generation neutrino detector will also be left for future work [113]. The stiffness of the hadronic phase also alters the resistance to PNS contraction, creating potential degeneracies with the onset properties of hybrid EOSs. A survey of hadronic EOSs with varying stiffness to quantify this degeneracy can assess whether AIC neutrino signals can disentangle EOS stiffness from PT effects.

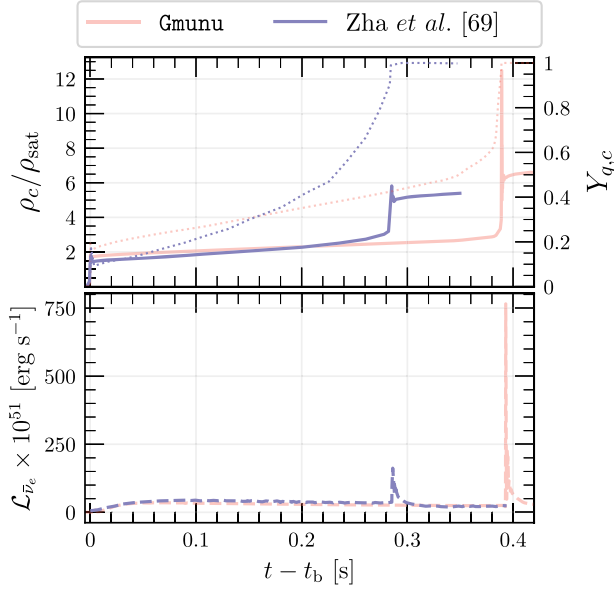


FIG. 6. Time evolution of central density ρ_c (solid), central quark fraction $Y_{q,c}$ (dotted) and luminosity of electron antineutrino $\mathcal{L}_{\bar{\nu}_e}$ (dashed). Results from Gmunu and Ref. [69] are colored in red and purple respectively.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX A: CORE COLLAPSE OF A $12M_\odot$ STAR WITH HYBRID EOS

In this section, we present general relativistic neutrino-radiation hydrodynamics simulations on the first-order

QCD PT in CCSNe using hybrid EOS as a reference result of QCD PT in Gmunu.

Our simulation is identical to that of the main text except for employing a different hybrid EOS, progenitor and computational domain. This hybrid EOS employs the Shen-Toki-Oyamatsu-Sumiyoshi EOS for hadronic phase [122] and MIT bag model with bag constant ($B = 165 \text{ MeV}$)⁴ for deconfined quark phase [123]. This EOS adopts Gibbs construction for the mixed nuclear-quark phase. This EOS is adopted from several studies of PT in CCSNe [67,69,124] and is publicly available [89]. We take the same $12M_\odot$, solar metallicity, presupernova progenitor $s12$ from [51], as in [69]. The computational domain covers from 0 to $2 \times 10^4 \text{ km}$, with the resolution $N_r = 128$ and allowing 11 AMR level. The finest grid size Δr at the star’s center is 153 m.

We compare our results with those from the axisymmetric simulation in [69], which employs an energy-independent neutrino transport scheme. Additionally, it solves the Newtonian Euler equations with general relativistic corrections. The computational domain extends from 0 to $2 \times 10^4 \text{ km}$ in the radial direction and spans $\pm 2 \times 10^4 \text{ km}$ along the cylindrical axis, with a finest spatial resolution of 150 m as reported in [69].

The time evolution of the central density ρ_c , central quark fraction $Y_{q,c}$, and electron antineutrino luminosity $\mathcal{L}_{\bar{\nu}_e}$ is shown in Fig. 6. The iron core collapses until the central density ρ_c exceeds nuclear saturation density ($\gtrsim 10^{14} \text{ g cm}^{-3}$), at which point the stiffening of the EOS halts the collapse and triggers a core bounce, followed by continued mass accretion. In this model, the PNS enters the mixed phase even before the core bounce, near ρ_{sat} , due to the low onset density in the hybrid EOS. With ongoing mass accretion, Kelvin-Helmholtz cooling, and EOS softening in the mixed phase, the PNS contracts rapidly, accompanied by a substantial rise in the central quark fraction $Y_{q,c}$. This is reflected in the significant increase in both ρ_c and $Y_{q,c}$ within approximately 0.2 s after the initial core bounce.

Here, we observe a difference in the rate of increase of ρ_c and $Y_{q,c}$, primarily due to differences in the neutrino transport schemes. Our simulation adopts an energy-averaged transport scheme, while Ref. [69] employs an energy-dependent transport approach.

We observe a second dynamical collapse and the formation of a PHS in both our simulation and Ref. [69], occurring when the central quark fraction reaches $Y_{q,c} \sim 0.3$. This transition leads to a strong collapse, with the central density reaching up to $\sim 12\rho_{\text{sat}}$ in our simulation and $\sim 6\rho_{\text{sat}}$ in [69]. The timing of the second collapse also differs: $t - t_b = 0.389 \text{ s}$ in our simulation compared to $t - t_b = 0.286 \text{ s}$ in [69], primarily due to differences in the neutrino transport scheme, gravitational treatment, and simulation dimensionality. Additionally, the luminosity of the second electron antineutrino burst differs by a factor of

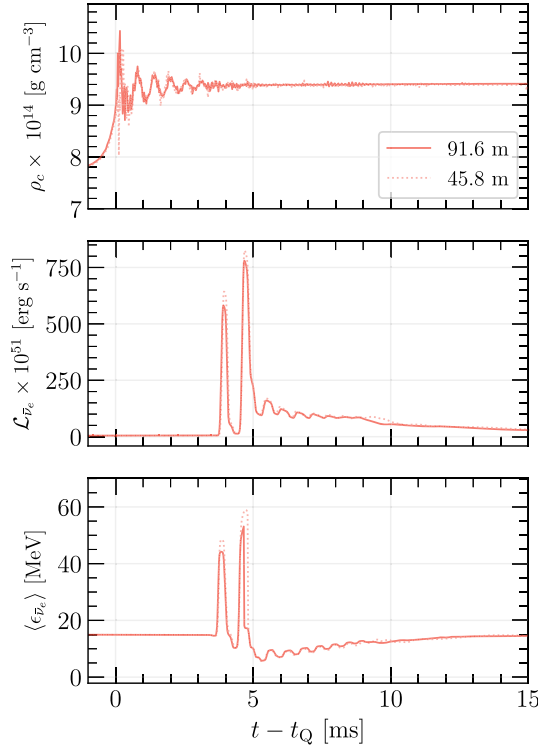


FIG. 7. Time evolution of the central density ρ_c , luminosity of electron antineutrino $\mathcal{L}_{\bar{\nu}_e}$ and corresponding average energy $\langle \epsilon_{\bar{\nu}_e} \rangle$ after the second core bounce for RDF-1.9 models with different resolution. The models with finest resolutions of 91.6 and 45.8 m are shown in solid and dotted lines, respectively.

approximately 3, driven by the same factors. Despite these quantitative differences arising from the numerical setups, our results remain qualitatively consistent with those of Ref. [69].

With major differences in the setup, this set of simulations does not constitute a systematic comparison of specific neutrino microphysics or numerical choices. However, the results highlight how numerical schemes and dimensionality can alter PT dynamics and the associated neutrino observables in CCSN. It is therefore reasonable to expect that PT in AIC systems will be subject to similar sensitivities. A key observable difference between AIC systems and CCSNe is that AIC systems exhibit a significantly longer time delay between the first and second neutrino bursts.

APPENDIX B: RESOLUTION DEPENDENCE

We perform a set of simulations with different resolution with the RDF-1.9 model to compare the main result discussed in the main text. The numerical setup is identical to that in the main text, except that 9 and 10 AMR levels are allowed, with a finest solution of 91.6 and 45.8 m.

The PNS cores of both models enter the mixed phase at $t_{\text{mixed}} = 0.58$ s and the deconfined quark phase at $t_Q = 1.1$ s. Figure 7 shows the evolution of ρ_c , $\mathcal{L}_{\bar{\nu}_e}$, and

$\langle \epsilon_{\bar{\nu}_e} \rangle$ around t_Q . The core oscillations agree qualitatively well in both frequency and amplitude across the two resolutions. Similarly, the amplitude and oscillatory features of the second neutrino burst are in good agreement between the two resolutions.

APPENDIX C: COMPARISON OF TWO NEUTRINO BURSTS IN AIC

Electron captures prior to core bounce steadily increase $\mathcal{L}_{\nu_e}$. The collapse halts due to hadronic matter stiffening, triggering an outward shock that dissociates heavy nuclei, enhancing ν_e production via e^- captures. This results in a luminous ν_e burst from proton captures in shock-heated matter, peaking at $t = t_b$ in Fig. 8. Compression and Doppler redshifting cause a dip in $\mathcal{L}_{\nu_e}$ around 4 ms later.

As the shock breaks out, $\mathcal{L}_{\text{tot}}$ peaks at 4.6×10^{53} erg s⁻¹, with $\mathcal{L}_{\nu_e}$ peaks at 3.8×10^{53} erg s⁻¹, while shock-heated matter enhances pair processes, leading to $\bar{\nu}_e$ and ν_x peaks at 4×10^{52} erg s⁻¹ and 6×10^{52} erg s⁻¹, respectively. Given the similarity of hadronic matter modeling near nuclear saturation density across RDF EOS models, this first neutrino burst behavior is generic.

In contrast, the second neutrino burst differs significantly due to a different emission mechanism. The stiffening of deconfined quark matter halts the PT-induced collapse, generating a second shock. The low-proton-fraction cavity facilitates e^+ capture, leading to a $\bar{\nu}_e$ -dominated burst. Unlike the first burst, all species reach a local maximum in Fig. 8. $\mathcal{L}_{\text{tot}}$ peaks at 1.6×10^{53} erg s⁻¹ with a $\mathcal{L}_{\nu_e}$, $\mathcal{L}_{\bar{\nu}_e}$ and $\mathcal{L}_{\nu_x}$ peak at 6×10^{52} erg s⁻¹, 4×10^{52} erg s⁻¹, 2×10^{52} erg s⁻¹ respectively. Multiple follow-up shocks from PHS oscillations introduce an oscillatory pattern.

The relative luminosities of the first and second bursts depend on EOS properties, particularly the onset density, which affects the shock strength. (See discussion of Fig. 5). Consequently, the total luminosity of the model with RDF-1.9 and RDF-1.2 is larger than the first neutrino burst, with $\mathcal{L}_{\text{tot}} = 2.99 \times 10^{54}$ erg s⁻¹ and 1.35×10^{54} erg s⁻¹ respectively. On the other hand, $\mathcal{L}_{\text{tot}}$ of the second neutrino burst in the RDF-1.5 and RDF-1.1 models are weaker than the first, with $\mathcal{L}_{\text{tot}} = 2 \times 10^{53}$ erg s⁻¹ and 5.47×10^{52} erg s⁻¹, respectively. Thus, no clear correlations exist between the total luminosities of the first and second bursts, as these are highly sensitive to EOS.

APPENDIX D: COMPARISON OF PHASE TRANSITION IN AIC AND CCSN

To investigate how a massive envelope in the progenitor influences the onset conditions and neutrino observables, we perform four CCSN simulations using the same progenitor and setup as in Appendix A, except that the computational domain extends from $[0, 1 \times 10^4]$ km, with a resolution of $N_r = 128$ and allowing 11 AMR levels.

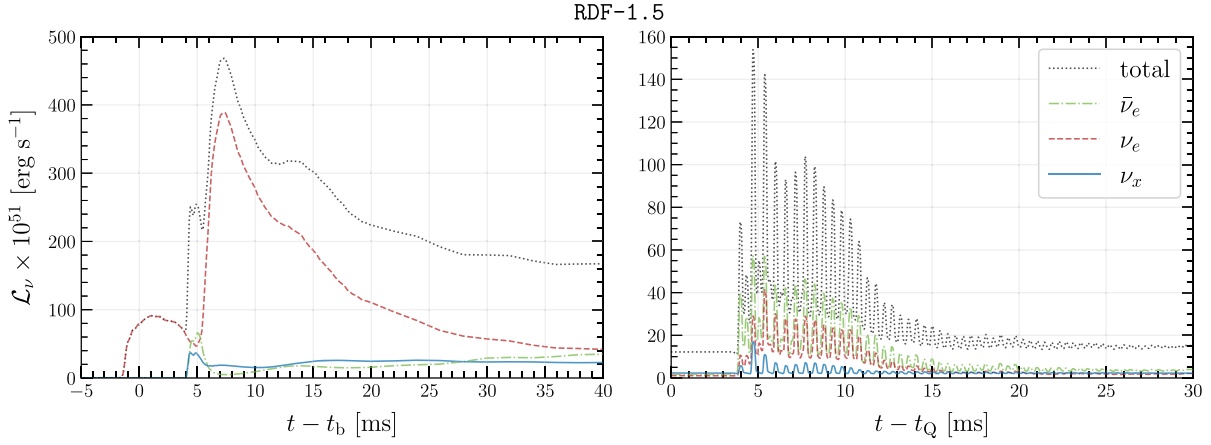


FIG. 8. Neutrino luminosity $\mathcal{L}_\nu$ of the first and the second burst in the model with RDF-1.5 for total luminosity (black dotted), $\mathcal{L}_{\bar{\nu}_e}$ (red dashed), $\mathcal{L}_{\nu_e}$ (green dash-dotted) and $\mathcal{L}_{\nu_x}$ (blue solid), respectively.

The finest grid size Δr at the center of the star is 76.3 m. The four RDF EOSs used are identical to those in the main text.

Similar to Appendix A, the collapse of the iron core is halted by the stiffening of the EOS when the central density ρ_c exceeds nuclear saturation density ($\gtrsim 10^{14} \text{ g cm}^{-3}$), leading to a core bounce. In contrast to the AIC case, the PNS enters the mixed phase only after core bounce due to the higher onset density characteristic of the RDF EOS family. The left columns of Table III present the time of mixed-phase entry t_{mixed} , central density $\rho_{\text{mixed,c}}$, central temperature $T_{\text{mixed,c}}$, the enclosed baryonic mass of the PNS M_{mixed}^b , and the total neutrino luminosity $\mathcal{L}_{\text{mixed,tot}}$ at the moment the system first enters the mixed phase, defined by $Y_q > 0$.

The baryonic mass at the onset of the mixed phase, M_{mixed}^b , varies by up to 15.8% across the CCSN models, indicating a prolonged accretion phase due to the presence of a massive envelope. In contrast, the AIC models exhibit a variation of less than 0.07% in M_{mixed}^b , reflecting the absence of extended mass accretion.

Due to the elevated central temperature T_c and low proton fraction Y_p in the PNS core, the modified onset density $\rho_{\text{mixed,c}}$ is lower than the cold, beta-equilibrated onset density ρ_{onset} listed in Table I. For instance, in model RDF-1.1, $\rho_{\text{mixed,c}}$ is 22% lower than ρ_{onset} , corresponding to $T_{\text{mixed,c}} = 27.5 \text{ MeV}$ (or 17.8 MeV) at $t = t_{\text{mixed}}$. This thermal reduction in the onset density is less pronounced in the CCSN system compared to the 30% reduction observed in the AIC system, as reported in the main text.

The presence of a massive envelope in CCSN leads to ongoing mass accretion and a faster contraction rate compared to AIC. As a result, PNSs in CCSN enter the mixed phase earlier than their AIC counterparts when evolved with the same RDF EOS, and generally exhibit similar or higher values of $\rho_{\text{mixed,c}}$. For example, the CCSN model with RDF-1.5 enters the mixed phase approximately 1 s earlier than in the corresponding AIC case, with a higher central density at the onset. Therefore, continuous accretion in CCSN can significantly influence the onset conditions for the QCD PT in the PNS.

TABLE III. Various quantities extracted at the moment of first reaching the mixed nuclear-quark phase (columns in the left part) in our CCSN simulations, using different RDF EOS. t_{mixed} , $\rho_{\text{mixed,c}}$, $T_{\text{mixed,c}}$, M_{mixed}^b and $\mathcal{L}_{\text{mixed,tot}}$ denote the time after the first core bounce at $t = t_b$, central density, central temperature, the enclosed baryon mass of the PNS and the total neutrino luminosity for the system first reaching the mixed nuclear-quark phase. t_Q and $M_{t_Q}^b$ is the time and the enclosed baryon mass of the PHS first reaching the deconfined quark phase with respect to t_b . The values of neutrino luminosity and averaged energy are evaluated at $r = 500 \text{ km}$. The values of the second neutrino burst are represented as $\mathcal{L}_{\nu_i}$ and $\langle \epsilon_{\nu_i} \rangle$, respectively for electron neutrino ν_e , electron antineutrino $\bar{\nu}_e$, and a collective species describing heavy lepton neutrinos and their antineutrinos as ν_x . Note that the PNS or PHS are defined as the regions of $\rho > 10^{11} \text{ g cm}^{-3}$.

EOS	$t_{\text{mixed}} [\text{s}]$	$\rho_{\text{mixed,c}} [10^{14} \text{ g cm}^{-3}]$	$T_{\text{mixed,c}} [\text{MeV}]$	$M_{\text{mixed}}^b [M_\odot]$	$\mathcal{L}_{\text{mixed,tot}} [10^{51} \text{ erg s}^{-1}]$	$t_Q [\text{s}]$	$M_{t_Q}^b [M_\odot]$	$\mathcal{L}_{\nu_e}$	$\mathcal{L}_{\bar{\nu}_e}$	$\mathcal{L}_{\nu_x}$	$\langle \epsilon_{\nu_e} \rangle$	$\langle \epsilon_{\bar{\nu}_e} \rangle$
RDF-1.9	0.4	4.1	14.3	1.58	79	0.81	1.66	363	914	66952.7	59.2	98.4
RDF-1.2	1.0	5.3	23.9	1.69	55.9	1.59	1.76	178	1234	333 40	61.4	90.4
RDF-1.5	1.9	6.26	22.9	1.78	39.8	2.27	1.8	188	752	372 29	42.2	112.4
RDF-1.1	2.7	6.84	24.7	1.81	23.3	2.94	1.82	183	875	33842.2	48	147.9

Upon the onset of the first-order PT, the contraction of the PNS accelerates due to EOS softening in the mixed phase and receives reduced support from Kelvin-Helmholtz cooling. This support is less pronounced at higher onset densities, as indicated by the decreasing trend in $\mathcal{L}_{\text{mixed,tot}}$. Although this trend is qualitatively similar to that observed in AIC, $\mathcal{L}_{\text{mixed,tot}}$ in CCSN is approximately a factor of 2 larger, primarily due to the larger M_{mixed}^b , and thus a greater gravitational binding energy. Consequently, the presence of a massive envelope in CCSN leads to a stronger neutrino cooling effect compared to AIC systems, resulting in a more rapid increase in central density.

With continued contraction and ongoing softening of the EOS, the PNS undergoes a supersonic collapse that is halted by the formation of a stiffer deconfined quark core. This results in the formation of a hydrodynamical shock, whose propagation generates a second neutrino burst primarily due to positron capture. This behavior is qualitatively consistent with the PT dynamics observed in the AIC models presented in the main text and is well documented in various studies of QCD PT in CCSNe [67,69–73,76].

The right columns in Table III present the onset time of the deconfined quark phase t_Q , the enclosed baryon mass of the PHS at the onset of the deconfined quark phase $M_{t_Q}^b$, the neutrino luminosity $\mathcal{L}_{\nu_l}$ and average energy $\langle \epsilon_{\nu_l} \rangle$ for all neutrino species associated with the second neutrino burst.

The secular mass accretion in CCSN continues throughout the PT, as indicated by the increase in the enclosed baryonic mass of the PNS ($M_{t_Q}^b > M_{\text{mixed}}^b$). In contrast, the change in baryonic mass during the PT in AIC is negligible.

This supports a more rapid progression of the PT in CCSN, as reflected by the shorter transition duration $t_Q - t_{\text{mixed}} \lesssim 0.6$ s across all models, compared to the longer timescale of up to ~ 3 s in the AIC case, such as for model RDF-1.1. These results highlight the significant role of a massive envelope in facilitating the PT in CCSN, making the evolution of the PNS less sensitive to the specific properties of the hybrid EOS described in the main text.

The onset condition of the deconfined quark phase determines the characteristics of the second hydrodynamical shock, and consequently, the associated neutrino observables. Since $M_{t_Q}^b$ varies among different CCSN models, an EOS with a higher onset density can lead to collapse in a more massive PNS, complicating the collapse dynamics. It is evident that, in CCSN, none of the neutrino observables exhibits a strictly consistent decreasing trend across all neutrino species. In contrast, both $\mathcal{L}_{\bar{\nu}_e}$ and $\langle \epsilon_{\nu_l} \rangle$ show a clear decreasing trend with increasing onset density in AIC models. Although a more detailed and systematic investigation of the PHS structure and collapse dynamics is required to fully understand how the onset condition influences the neutrino signal, it is clear that the presence of a massive envelope in CCSN can modify the onset conditions and alter the resulting neutrino observables. The presence of a massive envelope introduces a potential systematic uncertainty when using these observables to constrain the properties of the hybrid EOS. In contrast, such systematics are negligible in AIC due to the absence of a massive envelope, making AIC a cleaner environment for probing the EOS-dependent features of the PT.

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