

Nuclear liquid-gas phase transition and supernovae evolutionJérôme Margueron,^{1,*} Jesús Navarro,² and Patrick Blottiau³¹*GANIL CEA/DSM-CNRS/IN2P3 BP 5027 F-14076 Caen Cedex 5, France*²*IFIC (CSIC, Universidad de Valencia) Apdo. 22085, E-46.071-Valencia, Spain*³*CEA/DIF DPTA BP12 F-91680 Bruyères-le-Châtel Cedex, France*

(Received 8 April 2004; published 10 August 2004)

It is shown that the large density fluctuations appearing at the onset of the first order nuclear liquid-gas phase transition can play an important role in the supernovae evolution. Due to these fluctuations, the neutrino gas may be trapped inside a thin layer of matter near the protoneutron star surface. The resulting increase of pressure may induce strong particle ejection a few hundred milliseconds after the bounce of the collapse, contributing to the revival of the shock wave. The Hartree-Fock+RPA scheme, with a finite-range nucleon-nucleon effective interaction, is employed to estimate the effects of the neutrino trapping due to the strong density fluctuations, and to discuss qualitatively this self-consistent dynamical effect for supernovae evolution.

DOI: 10.1103/PhysRevC.70.028801

PACS number(s): 97.60.Bw, 26.50.+x, 25.30.Pt, 21.60.Jz

The first simulations of Colgate and White [1] and Arnett [2] have settled the general scenario of explosive supernovae, which schematically goes as follows. Stars with more than about ten times the Sun mass develop central iron cores which eventually become unstable and collapse to neutron stars, due to electron capture and photodisintegration onto iron-group nuclei. The interior of neutron stars is denser than nuclear matter and initially extremely hot. Above nuclear density, the equation of state becomes stiff enough to produce a bounce of the core, and a shock wave is formed, moving to the infalling outer core and envelope. Particle reactions at such conditions create a huge number of neutrinos, which eventually escape from the dense neutron star, transferring energy to the matter in the still infalling outer layers of the star. This neutrino heating is believed to cause the violent disruption of the star in a supernova explosion. However, from the very beginning it was found that the resulting shock wave is not energetic enough. Recently, an enormous effort has been done to describe the neutrino production and interactions in great detail [3], including the effects of stellar rotation and of violent anisotropic plasma motions. Convective processes in the supernova core had been recognized to accelerate the energy transport inside the neutron star and to enhance the deposition of energy by neutrinos in the outer stellar layers, thus supporting the explosion of the star (the “delayed” mechanism). Along some decades, numerical simulations were performed, starting from the previous 1D codes [4,5] and the flux-limited neutrino transport scheme [6] to reach the multidimensional calculations, showing the importance of convection and Rayleigh–Taylor instabilities, and the most refined neutrino transport methods [7,8]. Nevertheless, the outcome of these worldwide most elaborate supernova simulations is disappointing, as no explosions could be obtained. This negative result shatters the widely accepted view of how the explosion starts.

In this work we point out that large density fluctuations

associated to the first order nuclear liquid-gas phase transition, can play an important role in the neutrino trapping. Indeed, the scattering of neutrinos is hugely increased at the onset of strong density fluctuations, so that neutrinos should be trapped inside a thin layer of the protoneutron star (PNS) for densities in the range between about $0.1\rho_0$ and $0.6\rho_0$, where ρ_0 is the nuclear matter density at saturation. The resulting increase in the pressure may induce strong particle ejection a few hundred milliseconds after the bounce of the collapse¹, contributing to the revival of the shock wave.

A large nuclear community is currently studying the liquid-gas phase transition which occurs during multifragmentation experiments in heavy-ion collisions. The present status is reviewed in Ref. [9] and here we recall the main features directly related to our purpose. Such a transition appears below the saturation density ρ_0 of nuclear matter, at around $\sim 0.6\rho_0$, i.e., $\sim 0.1 \text{ fm}^{-3}$. Heavy nuclei are located on the liquefaction curve for $T=0 \text{ MeV}$ and the critical temperature is located between 15 and 20 MeV, depending on the specific nuclear model used. The onset of liquid-gas coexistence phase is associated to the spinodal instability, related to a negative curvature in the nucleonic free energy density such that the speed of sound vanishes. The analysis of the topology of the thermodynamic potential show that the phase transition is first order in symmetric and asymmetric nuclear matter. A unique spinodal instability is expected and the asymmetry do not modify its properties up to very large asymmetries [10,11]. The spinodal instability is isoscalar because the interaction between protons and neutrons is attractive at low density. Hence, in the density regime close to the spinodal, the density response function will be enhanced, thus affecting the neutral current neutrino opacity while the charge current neutrino opacity is not modified. This may result in a reduced total neutrino mean-free path.

It is worth pointing out that Sawyer [12] obtained a quite similar result investigating the dependence of the neutrino

*Present address: Institut für Theoretische Physik, Universität Tübingen, D-72076 Tübingen, Germany.

¹Authors estimation is based on the typical time neutrinos need to diffuse inside the protoneutron star (few tens of milliseconds) and the typical time to escape from it (few seconds).

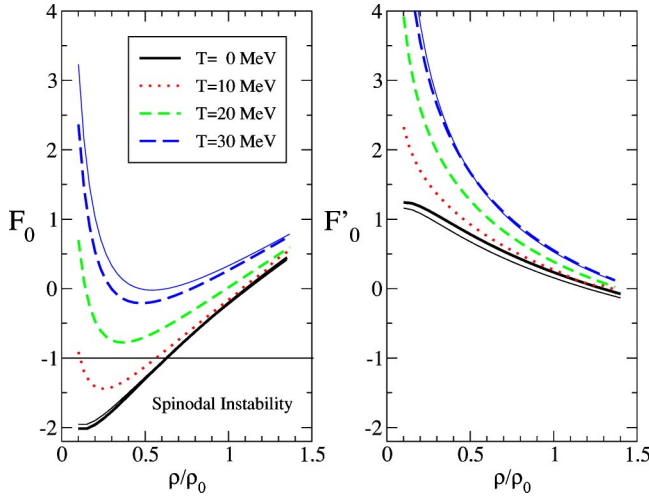


FIG. 1. (Color online) The adimensional Landau parameters F_0 and F'_0 are plotted as a function of the density at several values of temperature. The thin line represents a calculation of the same quantity in asymmetric nuclear matter with $x_p=0.25$ and for $T=0$ and 30 MeV. The definitions $F_0=(F_0^{nn}+F_0^{pp}+2F_0^{np})/2$ and $F'_0=(F_0^{nn}+F_0^{pp}-2F_0^{np})/2$ has been taken and the coefficients $F_0^{rr'}$ are the second derivative of the free energy density with respect to the densities. The onset for the spinodal instability is given by $F_0=-1$.

opacity on the equation of state of pure neutron matter. Assuming that the neutrino scattering is dominated by the classical fluctuations it was shown that the neutrino mean-free path is inversely proportional to the neutron matter compressibility. Consequently, a local softening of the equation of state could produce a thin layer, near to the star surface, quite opaque to neutrinos of some energy. In this work we consider the case of asymmetric nuclear matter, where a liquid-gas transition can be present at low densities, contrarily to the pure neutron case. We calculate consistently both the nuclear equation of state and the density fluctuations using the Gogny DIP finite-range nucleon-nucleon effective interaction [13]. We describe the nuclear isoscalar and isovector density fluctuations within the Hartree-Fock (HF) plus the random phase approximation (RPA). The calculation of the RPA response functions is performed in the Landau approximation, ignoring the Landau coefficients beyond the dipole one in both density and spin channels [14]. It is worth recalling that some values of Landau parameters are related to different instabilities. In particular, the density at which the dimensionless Landau parameter F_0 equals -1 indicates the onset of the spinodal instability. Figure 1 shows the Landau parameters F_0 and F'_0 versus the total density for symmetric nuclear matter and several densities. The thin line correspond to asymmetric nuclear matter with $x_p \equiv Z/(N+Z)=0.25$. The results in asymmetric nuclear matter look very similar and we clearly see that out of the spinodal region, F_0 is attractive for a wide range of densities and temperatures. The Landau parameter F'_0 which is responsible for isospin-flip fluctuations do not show any instable behavior, which confirms that the instability is isoscalar.

In Fig. 2 we plot the pressures and particle fractions as a function of the density for asymmetric nuclear matter at a temperature of $T=10$ MeV and leptonic fraction $Y_l=0.285$,

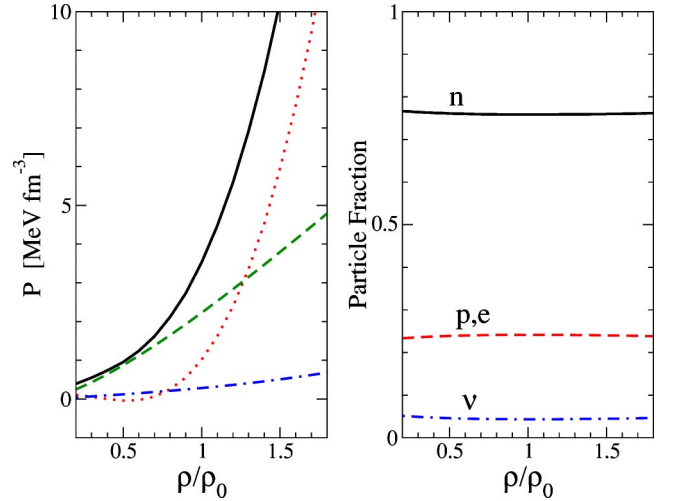


FIG. 2. (Color online) Pressures (left-hand panel) and particle fractions (right-hand panel) plotted as a function of the density for asymmetric nuclear matter at temperature $T=10$ MeV and leptonic fraction $Y_l=0.285$, calculated with the Gogny DIP effective interaction. The total pressure (solid line) is the sum of the nuclear (dotted line), the electronic (dashed line), and the neutrino (dot-dashed line) pressures. On the right-hand panel, solid, dashed, and dot-dashed lines refer to the particle fraction of neutrons, charged particles (i.e., protons and electrons), and neutrinos, respectively.

defined as $(\rho_e + \rho_\nu)/\rho$. In the low density region $\rho < \rho_0$, the total pressure is mainly supported by electrons, and the contribution of neutrinos is about 20% or less [15].

In the nonrelativistic limit, the neutrino mean-free path λ is given by the well-known expression [16]

$$\frac{1}{\lambda(k_\nu, T)} = \frac{G_F^2}{16\pi^2} \int d\mathbf{k}_3 [c_V^2 (1 + \cos \theta) S^{(0)}(q, T) + c_A^2 (3 - \cos \theta) S^{(1)}(q, T)], \quad (1)$$

where k_ν is the neutrino energy-momentum, T is the temperature, G_F is the Fermi constant, $c_V(c_A)$ the vector (axial) coupling constant, \mathbf{k}_3 the final neutrino momenta, $q = k_\nu - k_3$ the transferred energy momentum, and $\cos \theta = \hat{\mathbf{k}}_\nu \cdot \hat{\mathbf{k}}_3$. The dynamical structure factors $S^{(S)}(q, T)$ describe the response of nuclear matter to excitations induced by neutrinos, and they contain the relevant information on the medium. The vector (axial) part of the neutral current gives rise to density (spin-density) fluctuations, corresponding to the $S=0$ ($S=1$) spin channel. Density and spin-density responses contribute to the scattering of neutrinos, while the isospin-flip response contribute to neutrino absorption process, which is known to dominate the total cross section [17]. We shall show that large isoscalar fluctuations can give the dominant contribution at subnuclear densities.

Clearly, an enhancement of the response function in either channel will result in a quenching of the neutrino mean-free path. It has been shown [18] that at densities higher than ρ_0 a ferromagnetic instability can appear, leading consequently to a zero neutrino mean-free path [19]. However, the density at which such an instability could appear is not yet a settled

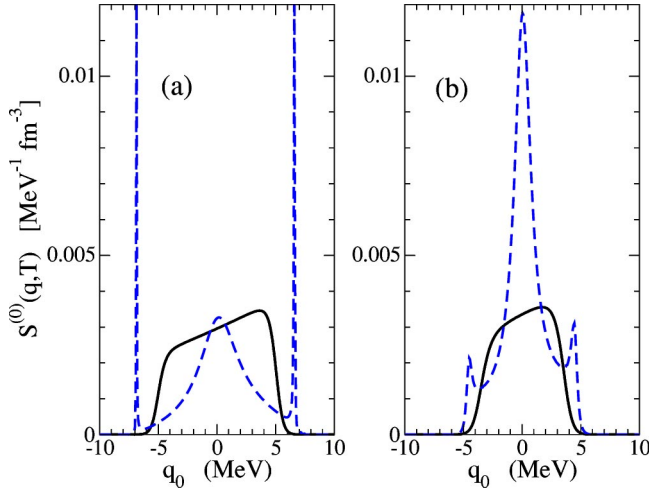


FIG. 3. (Color online) The density-density response function $S^{(0)}(q, T)$ is plotted as a function of the transferred energy for two values of the density: (a) ρ_0 and (b) $0.75\rho_0$. The calculations have been done for momentum $q=10$ MeV, temperature $T=10$ MeV and leptonic fraction $Y_l=0.285$. The solid and dashed lines stand for the HF and the RPA results, respectively.

question. For instance, whereas all Skyrme effective interactions predict a ferromagnetic transition for densities less than $\sim 3.5\rho_0$ [20], recent microscopic calculations [21,22] push it to higher densities, in case it exists.

Now we concentrate on the mechanism responsible for neutrino trapping. In the following, we will approach the liquid-gas phase transition from the homogeneous nuclear matter. This can be done by decreasing either the density at constant temperature or the temperature at constant density. In both cases, in order to avoid the phase coexistence region, we should keep the density beyond the spinodal one or the temperature beyond the critical one. The density-density response function $S^{(0)}(q, T)$ is represented in Fig. 3 for density values approaching the spinodal point. The solid line stands for the HF calculation while the dashed line includes the RPA correlations, with a particle-hole interaction. In the left-hand panel, we see two symmetric collective states inside the continuum strength. These states indicate that near the Fermi surface, the residual interaction is repulsive. It can be seen that when the system approaches the liquid-gas phase transition, a zero energy mode becomes increasingly enhanced. In fact, the peak diverges at the spinodal density. This is a signature of a strong attractive residual interaction, which indeed induces the phase transition.

The neutrino mean free path as a function of the density is plotted in Fig. 4 for typical thermodynamical quantities. The HF results are strongly dependent on the neutrino energy, as it has been shown by several authors [12,16,17], due mainly to the Pauli kinematical factors. One can see that the RPA correlations induce a strong reduction of the neutrino mean-free path close to the liquid-gas phase transition. This trapping is independent of the neutrino energy. Hence, neutrinos could be massively trapped inside the thin external layer with densities near the densities of the liquid-gas phase transition, and a radius close to the radius of the protoneutron star. The spinodal density is close to $0.6\rho_0$ in symmetric nuclear mat-

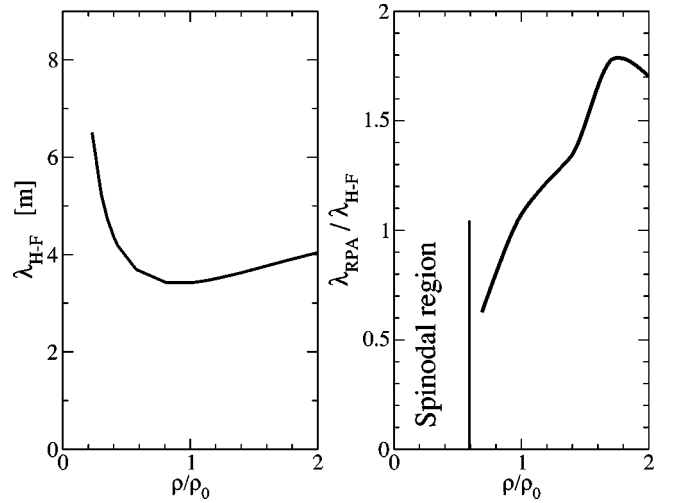


FIG. 4. The neutrino mean-free path is plotted as a function of the density for asymmetric nuclear matter at $T=10$ MeV and $Y_l=0.285$, the neutrino energy being equal to the neutrino chemical potential. Left panel: HF result, right panel: ratio between RPA and HF mean-free paths.

ter and is nearly constant in asymmetric nuclear matter until $\rho_x/\rho \sim 0.1$ ($x=n$ or p) [11].

Let us assume the extreme hypothesis that, after a few hundred milliseconds, all the neutrinos contained inside the star are trapped inside the thin layer near the surface, and that the initial distribution of neutrino density is quasiuniform. Therefore, the increase of the neutrino pressure inside the thin layer is given by the ratio of the PNS radius R over the width ϵ of the thin layer and writes

$$\frac{P_{\text{layer}}}{P_{\text{inside}}} \sim \left(\frac{R}{3\epsilon}\right)^{4/3}, \quad (2)$$

where the neutrino pressure is polytropic versus the density with the power $4/3$. Hence, for a typical radius of 10 km and a width of 100 m, the neutrino pressure increases by a factor greater than 100. Of course, the Pauli principle should strongly moderate this factor, but this schematic calculation points out the possible efficiency of the neutrino trapping.

This qualitative calculation assumes that all species are independent from each other. In fact, neutrinos are in β equilibrium, so that a modification of their number implies a corresponding change of the number of other particles, in particular the electrons. If the neutrino density is increased at constant temperature and constant baryonic density, then the β equilibrium shifts in the direction to raise the number of electrons and protons. This is the consequence of Le Chatelier's principle of minimum constraint extended to quantum systems. The same principle predicts that when the temperature is increased, the system tends to move towards the endothermic direction, thus increasing the number of electrons and protons. Consequently, raising both neutrino density and temperature produces electrons and protons. As electrons are mainly responsible for the pressure at low density, the overall effect of neutrino trapping is the increase of the total pressure.

Let us now analyze the physical consequences for the supernovae-II mechanism due to the interplay between neutrinos propagation and the liquid-gas phase transition. During the iron core collapse, neutrinos are trapped inside the core. The energy gained by the collapse is essentially carried away by the neutrinos after the bounce. Neutrinos are massively trapped for densities below the saturation one, and hence those neutrinos which are diffusing from the inner part of the PNS cannot reach the outer part of the star. An important fraction of neutrinos accumulates inside a thin layer close to the surface of the PNS. The neutrino Fermi energy increases and shifts the beta equilibrium to the production of protons and electrons. The neutrino and electronic pressures increase as well as the temperature inside the thin layer. We believe that this dynamical effect may lead to a partial conversion of the energy carried away by the neutrino gas to the matter through the beta equilibrium. According to the Euler transport equation, a local increase of pressure will produce an ejection of matter. This radiated matter will encounter the shock wave and may contribute to its revival. Indeed, the energy carried away by neutrinos represents about 98% of the collapse energy while the shock wave transports only 2% of it. If the ejection of particles carries away only 2% of neutrino energy, the total kinetic energy of matter will increase by a factor of 2. This is the reason why we believe that this new scenario could contain an efficient mechanism for the explosion.

In this work we have explored the consequences that large density fluctuations close to the liquid-gas phase transition could have for neutrino propagation and supernovae-II explosion. The calculations presented here should apply only to the homogeneous phase close to the coexistence region, either for densities higher than the spinodal one, or temperatures higher than the critical one. For instance, close to the critical temperature, we expect to observe the coherent scattering of neutrinos in analogy with the well-known critical opalescence phenomenon (with respect to photons) observed in condensed matter. Recently, neutrino opacity has also been found to become very high inside the pasta phase, due to the coherent scattering of neutrinos off clusters [23]. This indicates that inhomogeneous nuclear matter should also trap neutrinos, and increase the efficiency of our scenario. We propose to include the trapping in hydrodynamical simulations including the self-consistent shift of the β equilibrium. The dynamics should result in a bump of the pressure close to the surface of the PNS. Work in this direction is now in progress.

Discussions with M. Baldo, Ph. Chomaz, H. Mütter, A. Pérez-Cañellas, and A. Sedrakian are gratefully acknowledged. J.N. is supported by MCyT/FEDER and Generalitat Valenciana (Spain), Grant Nos. BMF2001-0262 and GV2003-002.

-
- [1] S. A. Colgate and R. H. White, *Astrophys. J.* **143**, 626 (1966).
 - [2] W. D. Arnett, *Can. J. Phys.* **44**, 2553 (1966).
 - [3] R. Buras, M. Rampp, H.-Th. Janka, and K. Kifonidis, *Phys. Rev. Lett.* **90**, 241101 (2003).
 - [4] H. A. Bethe and J. R. Wilson, *Astrophys. J.* **295**, 14 (1985).
 - [5] P. Blottiau, Ph.D. thesis, Paris VII University, 1989.
 - [6] S. Bruenn, *Astrophys. J., Suppl. Ser.* **58**, 771 (1985).
 - [7] P. Mellor, J. P. Chièze, and J. L. Basdevant, *Astron. Astrophys.* **197**, 123 (1988).
 - [8] M. Rampp and H. T. Janka, *Astron. Astrophys.* **396**, 361 (2002).
 - [9] Ph. Chomaz, M. Colonna, and J. Randrup, *Phys. Rep.* **389**, 263 (2004).
 - [10] V. Baran, M. Colonna, M. Di Toro, and V. Greco, *Phys. Rev. Lett.* **86**, 4492 (2001).
 - [11] J. Margueron and Ph. Chomaz, *Phys. Rev. C* **67**, 041602(R) (2003).
 - [12] R. Sawyer, *Phys. Rev. D* **11**, 2740 (1975).
 - [13] M. Farine, D. Van-Eiff, P. Schuck, J.-F. Berger, J. Dechargé, and M. Girod, *J. Phys. G* **25**, 863 (1999).
 - [14] J. Margueron, Ph.D thesis, Paris XI University, 2001.
 - [15] E. Suraud, Ph.D thesis, Paris VI University, 1984.
 - [16] N. Iwamoto and C. J. Pethick, *Phys. Rev. D* **25**, 313 (1982).
 - [17] S. Reddy, M. Prakash, and J. M. Lattimer, *Phys. Rev. D* **58**, 013009 (1998).
 - [18] A. Vidaurre, J. Navarro, and J. Bernabeu, *Astron. Astrophys.* **135**, 361 (1984).
 - [19] J. Navarro, E. S. Hernández, and D. Vautherin, *Phys. Rev. C* **60**, 045801 (1999).
 - [20] J. Margueron, J. Navarro, and N. V. Giai, *Phys. Rev. C* **66**, 014303 (2002).
 - [21] S. Fantoni, A. Sarsa, and K. E. Schmidt, *Phys. Rev. Lett.* **87**, 181101 (2001).
 - [22] I. Vidaña and I. Bombaci, *Phys. Rev. C* **66**, 045801 (2002).
 - [23] C. J. Horowitz, M. A. Pérez-García, and J. Piekarewicz, *Phys. Rev. C* **69**, 045804 (2004).