NEUTRINOS FROM PROTONEUTRON STARS: 
A PROBE OF HOT AND DENSE MATTER

SANJAY REDDY AND MADAPPA PRAKASH

Physics Department
State University of New York at Stony Brook
Stony Brook, NY 11794-3800, USA

ABSTRACT

Neutrino processes in dense matter play a key role in the dynamics, deleptonization and the early cooling of hot protoneutron stars formed in the gravitational collapse of massive stars. Here we calculate neutrino mean free paths from neutrino-hyperon interactions in dense matter containing hyperons. Significant contributions to the neutrino opacity arise from scattering involving the $\Sigma^-$ hyperon, and absorption processes involving the neutral and $\Sigma^-$ hyperons. The estimates given here emphasize the need for (a) opacities which incorporate many-body effects in a multi-component mixture, and (b) new calculations of thermal and leptonic evolution of protoneutron stars with neutrino transport and equations of state with strangeness-rich matter.

1. Introduction

The general nature of the neutrino signature expected from a newly formed neutron star (hereafter referred to as a protoneutron star) has been theoretically predicted and confirmed by the observations from supernova SN1987A. Although neutrinos interact weakly with matter, the high baryon densities and neutrino energies achieved after the gravitational collapse of a massive star ($\geq 8$ solar masses) cause the neutrinos to become trapped on the dynamical timescales of collapse\textsuperscript{3,4}. Trapped neutrinos at the star’s core have Fermi energies $E_\nu \sim 200 - 300$ MeV and are primarily of the $\nu_e$ type. They escape after diffusing through the star exchanging energy with the ambient matter, which has an entropy per baryon of order unity in units of Boltzmann’s constant. Eventually they emerge from the star with an average energy $\sim 10 - 20$ MeV and in nearly equal abundances of all three flavors, both particle and anti-particle.

Although the composition and the equation of state of the hot protoneutron star matter are not yet known with certainty, QCD based effective Lagrangians have opened up intriguing possibilities. Among these is the possible existence of matter with strangeness to baryon ratio of order unity. Strangeness may be precipitated either in the form of fermions, notably the $\Lambda$ and $\Sigma^-$ hyperons, or, in the form of a Bose condensate, such as a $K^-$-meson condensate (see Ref.\textsuperscript{5} for detailed discussion
and extensive references). In the absence of trapped neutrinos, strange particles are expected to appear around $2 - 4$ times the nuclear matter density of $n_0 = 0.16 \text{ fm}^{-3}$. Neutrino-trapping causes the strange particles to appear at somewhat higher densities than in neutrino-free matter. The compositions shown in Figs. 1 and 2 highlight the influence of hyperons in the neutrino trapped regime. The results shown in these figures were calculated using a field-theoretical model in which baryons interact via the exchange of $\sigma$, $\omega$ and $\rho$ mesons. With the appearance of hyperons in matter, the electron-neutrino fraction increases with density in contrast to the case in which matter contains nucleons and leptons only. A similar behavior is observed in kaon condensed matter and also in matter where a phase transition to quark matter occurs. This behavior is associated with the presence of non-leptonic negatively charged particles in matter, such as the $\Sigma^-$ hyperon, or $K^-$ meson, or $d$ and $s$ quarks.

Keil and Janka have recently investigated the influence of the equation of state on the cooling and evolution of the protoneutron star. They find that the neutrino luminosity depends sensitively on the composition and on the stiffness of the equation of state at high density. In particular, the influence of hyperons, which introduces a softening of the high density equation of state, was examined. In many cases, the protoneutron star collapsed to a black hole causing an abrupt cessation of the neutrino signal. Although a clear distinction between the different equations of state could be achieved on the basis of the calculated neutrino signals, Keil and Janka conclude that “none of the models could be considered as a good fit of the neutron star formed in SN 1987A”.

We note, however, that in these studies, the sensitivity of the neutrino signals due only to the structural changes caused by the equation of state was assessed. In the presence of hyperons, the transport of neutrinos is also affected due to the changes in the composition, and additionally, to the interactions of neutrinos with strange particles. These effects were ignored in Ref. 8. Previous work involving neutrino interactions with hyperons was concerned with charged-current reactions only. Studies of the opacities and the transport processes of neutrinos out of the core containing nucleons, leptons, and in some cases, pion condensates may be found in Refs. 11 through 26.

It is our purpose here to study neutrino mean free paths in matter containing hyperons. Specifically, we will estimate scattering and absorption mean free paths of neutrinos when matter is under degenerate conditions. We find that significant contributions to the neutrino opacity arise from scattering involving the $\Sigma^-$ hyperon, and absorption processes involving the neutral and $\Sigma^-$ hyperons. Compared with ordinary matter, the presence of strangeness in matter is expected to lead to an excess of $\nu_e$ neutrinos relative to other types. Calculations of the complete thermal and leptonic evolution of a newly formed neutron star incorporating neutrino transport and equations of state with strangeness-rich matter will be taken up separately. With new generation neutrino detectors capable of recording thousands of neutrino events, it may be possible to distinguish between different scenarios observationally.
Fig. 1: The composition, electron chemical potential and entropy per particle in nucleons only matter with a lepton fraction $Y_{Le} = Y_e + Y_{\nu_e} = 0.4$, where $Y_i = n_i/n_b$. Results are from Ref. [5].
Fig. 2: The composition, lepton chemical potentials and entropy per particle in strangeness-rich matter with a lepton fraction $Y_{Le} = Y_e + Y_{\nu e} = 0.4$, where $Y_i = n_i/n_b$. Results are from Ref. [5].
2. Neutrino interactions with strange baryons

Neutrino interactions with matter proceed via charged and neutral current reactions. The neutral current processes contribute to elastic scattering, and the charged current reactions result in neutrino absorption. The interaction Lagrangian for these reactions is given by the Weinberg-Salam theory:

\[ \mathcal{L}_{\text{int}}^{nc} = \frac{G_F}{\sqrt{2}} l_\mu j_\mu^\nu \quad \text{for} \quad \nu + B \rightarrow \nu + B \]

\[ \mathcal{L}_{\text{int}}^{cc} = \frac{G_F}{\sqrt{2}} l_\mu j_\mu^\nu \quad \text{for} \quad \nu + B_1 \rightarrow \ell^- + B_2, \quad (1) \]

where \( G_F \simeq 1.436 \times 10^{-49} \text{ erg cm}^{-3} \) is the weak coupling constant, \( \nu \) is a neutrino, \( B_1 \) and \( B_2 \) are baryons, and \( \ell^- \) is a lepton. The leptonic and hadronic currents appearing above are:

\[ l_\mu = \bar{\psi}_\nu \gamma_\mu (1 - \gamma_5) \psi_\nu \]

\[ j_\mu^\nu = \bar{\psi}_e \gamma_\mu (1 - \gamma_5) \psi_\nu \]

\[ j_\mu^\nu = \bar{\psi}_i \gamma_\mu (g_{Vi} - g_{Ai} \gamma_5) \psi_i \]

\[ j_\mu^\nu = \bar{\psi}_i \gamma_\mu (C_{Vi} - C_{Ai} \gamma_5) \psi_i, \quad (2) \]

where \( i = n, p, \Lambda, \ldots \). The neutral current process couples neutrinos of all types (\( e, \mu \) and \( \tau \)) to the weak neutral hadronic current \( j_\mu^\nu \). The charged current processes of interest here are electron and muon neutrinos coupled to the charged hadronic current \( j_\mu^\nu \). The vector and axial vector coupling constants are listed in Table 1. Numerical values of the parameters that best fit the experiments are: \( D=0.756 \), \( F=0.477 \), \( \sin^2 \theta_W=0.23 \) and \( \sin \theta_c = 0.231 \).

### Table 1

| Reaction                | \( C_V \) | \( C_A \) | \( R_{nc}^{(1)} \) | \( R_{nc}^{(2)} \) |
|-------------------------|-----------|-----------|----------------|----------------|
| \( \nu_i + n \rightarrow \nu_i + n \) | -1        | -D - F   | 1          | 1           |
| \( \nu_i + p \rightarrow \nu_i + p \) | \( 1 - 4 \sin^2 \theta_W \) | \( D + F \) | 0.7504       | 0.8597       |
| \( \nu_i + \Lambda \rightarrow \nu_i + \Lambda \) | 0         | 0         | 0          | 0           |
| \( \nu_i + \Sigma^- \rightarrow \nu_i + \Sigma^- \) | \( -2 + 4 \sin^2 \theta_W \) | -2F       | 0.7392      | 0.6788       |
| \( \nu_i + \Sigma^+ \rightarrow \nu_i + \Sigma^+ \) | \( 2 - 4 \sin^2 \theta_W \) | 2F       | 0.7392      | 0.6788       |
| \( \nu_i + \Sigma^0 \rightarrow \nu_i + \Sigma^0 \) | 0         | 0         | 0          | 0           |
| \( \nu_i + \Xi^- \rightarrow \nu_i + \Xi^- \) | \( -1 + 4 \sin^2 \theta_W \) | D         | 0.2845     | 0.3238       |
| \( \nu_i + \Xi^0 \rightarrow \nu_i + \Xi^0 \) | 1         | -D - F   | 0.2681     | 0.1852       |

**NOTE.** The quantity \( R_{nc}^{(1)} = \frac{C_V^2 + 2C_A^2}{|V_{\nu B}|^2} / \left[ \frac{C_V^2 + 2C_A^2}{|V_{\nu n}|^2} \right] \) and \( R_{nc}^{(2)} = \frac{C_V^2 + 4C_A^2}{|V_{\nu B}|^2} / \left[ \frac{C_V^2 + 4C_A^2}{|V_{\nu n}|^2} \right] \).
TABLE 2
CHARGED CURRENT VECTOR AND AXIAL COUPLINGS

| Reaction                        | $g_V$   | $g_A$   | $R_{cc}$ |
|---------------------------------|---------|---------|----------|
| $\nu_i + n \rightarrow e + p$  | 1       | $F + D$ | 1        |
| $\nu_i + \Lambda \rightarrow e + p$ | $-\sqrt{3/2}$ | $-\sqrt{3/2}(F + D/3)$ | 0.0394 |
| $\nu_i + \Sigma^- \rightarrow e + n$ | $-1$     | $-(F - D)$ | 0.0125 |
| $\nu_i + \Sigma^- \rightarrow e + \Lambda$ | 0       | $\sqrt{2}/3$ | 0.2055 |
| $\nu_i + \Sigma^- \rightarrow e + \Sigma^0$ | $\sqrt{2}$ | $\sqrt{2}F$ | 0.6052 |
| $\nu_i + \Xi^- \rightarrow e + \Lambda$ | $\sqrt{3/2}$ | $\sqrt{3/2}(F - D/3)$ | 0.0175 |
| $\nu_i + \Xi^- \rightarrow e + \Xi^0$ | $\sqrt{1/2}$ | $(F + D)/\sqrt{2}$ | 0.0282 |
| $\nu_i + \Xi^0 \rightarrow e + \Sigma^+$ | 1       | $F + D$ | 0.0564 |
| $\nu_i + \Xi^0 \rightarrow e + \Sigma^+$ | 1       | $F - D$ | 0.2218 |

NOTE.– The quantity $R_{cc} = \left[ C^2(g_V^2 + 3g_A^2) \right]_{\nu + B_1 \rightarrow \ell + B_2} / \left[ C^2(g_V^2 + 3g_A^2) \right]_{\nu + n \rightarrow \ell + p}$.

In what follows, we consider the lowest order (tree level) processes for both elastic and absorption reactions. The squared matrix element for neutral current reactions is given by

\[ |M_{12\rightarrow34}|^2 = 16G_F^2 \left[ (C_V + C_A)^2 \left( p_1 \cdot p_2 \right) \left( p_3 \cdot p_4 \right) \right. \]
\[ \left. + (C_V - C_A)^2 \left( p_1 \cdot p_4 \right) \left( p_2 \cdot p_3 \right) \right. \]
\[ \left. - (C_V^2 - C_A^2) \left( p_2 \cdot p_4 \right) \left( p_1 \cdot p_3 \right) \right] , \]  

(3)

where the overline on $M$ denotes a sum over final spins and an average over the initial spins, and $p_i$ denotes the four momenta of particles $i = 1, 4$. The squared matrix element for the charged current reactions is given by a similar relation, but with the replacement $C_V \rightarrow g_V$, $C_A \rightarrow g_A$, and $G_F \rightarrow G_FC$, where $C = \cos \theta_C$ for a change of strangeness $\Delta S = 0$ and $C = \sin \theta_C$ for $\Delta S = 1$, consistent with the Cabibbo theory.

3. Neutrino mean free paths in degenerate matter

We turn now to consider the mean free path of neutrinos in stellar matter comprised of degenerate baryons (neutrons, protons and hyperons) and leptons under conditions of charge neutrality and chemical equilibrium. For the estimates below, we employ the non-relativistic approximation for the baryons, so that the squared matrix element takes a simple form. We treat the neutrinos only in the degenerate and non-degenerate limits. Results for arbitrary neutrino degeneracy and with the full matrix element in Eq. (3) will be reported elsewhere.
For elastic collisions \(1 + 2 \rightarrow 3 + 4\), where \(1(3)\) denotes the initial (final) neutrino and \(2(4)\) denotes the initial (final) baryon \(B\), the scattering relaxation time may be calculated by linearizing the Boltzmann equation. For small departure from equilibrium, the inverse relaxation times from the various components are additive; thus

\[
\frac{1}{\tau_s} = \sum g_2 \int \prod_{i=2}^{4} \frac{d^3 p_i}{(2\pi)^3} \mathcal{W}_{fi} \left[ n_2(1 - n_3)(1 - n_4) - (1 - n_2)n_3n_4 \right].
\]

Above, the sum is over all species of baryons, \(g_2\) is their degeneracy, \(n_i\) are the equilibrium Fermi-Dirac distributions, and \(\mathcal{W}_{fi}\) is the scattering rate.

\[
\mathcal{W}_{fi} = \left( \prod_{i=1}^{4} 2E_i \right)^{-1} \frac{1}{(2\pi)^4} \delta^4(p_1 + p_2 - p_3 - p_4) |M_{12\rightarrow34}|^2.
\]

The relaxation time \(\tau_s\) characterizes the rate of change of the distribution function \(n_i\) due to interactions with species 2, and may be used to define a scattering mean free path \(\lambda_s = c\tau_s\).

For elastic scattering on heavy fermions, the momentum transfer is small. Thus, the scattering rate may be expressed as a function of the neutrino energy \(E_\nu\) and the neutrino scattering angle \(\theta\). In degenerate matter, where the participant particles lie on their respective Fermi surfaces, the phase space integration can be separated into angle and energy integrals. Thus for degenerate neutrinos and when \(k_BT \ll E_\nu v_{Fi}/c\), where \(v_{Fi}\) is the velocity at the Fermi surface of species \(i\), the inverse relaxation time is given by (see Ref. [19, 11] for the result in a single component system)

\[
\frac{1}{\tau_s} = \sum_i \frac{G_F^2}{12\pi^3} \left( C_{V1}^2 + 2C_{A1}^2 \right) m_B^2 (k_BT)^2 E_\nu \left[ \frac{\pi^2}{4} + \frac{(E_\nu - \mu_\nu)^2}{(kT)^2} \right],
\]

where the sum is over the baryonic components present in the system.

For non-degenerate neutrinos and when \(k_BT \ll E_\nu v_{Fi}/c\), the result of Ref. [19] may be generalized to give

\[
\frac{1}{\tau_s} = \sum_i \frac{G_F^2}{15\pi^3} \left( C_{V1}^2 + 4C_{A1}^2 \right) p_{Fi}^2 E_\nu^3.
\]

The relaxation time for absorption through charged current reactions can be calculated in a similar fashion. When neutrinos are degenerate, absorption on neutrons and similarly on hyperons, is kinematically allowed. In this case, the inverse absorption length is given by

\[
\frac{1}{\tau_a} = \sum_j \frac{G_F^2 C_{ij} C^2}{4\pi^3} (g_{Vj}^2 + 3g_{Aj}^2) m_B m_B (k_BT)^2 \mu_e \left[ \frac{\pi^2}{4} + \frac{(E_\nu - \mu_\nu)^2}{(kT)^2} \right] \Xi
\]

with \(\Xi = \theta(p_{B2} + p_e - p_{B1} - p_\nu) + \frac{p_{B2} + p_e - p_{B1} + p_\nu}{2E_\nu} \theta(p_\nu - |p_{B2} + p_e - p_{B1}|)\).
where \( \theta(x) = 1 \) for \( x \geq 1 \) and zero otherwise.

When neutrinos are non-degenerate and when absorption is kinematically allowed, the relaxation time is given by

\[
\frac{1}{\tau_a} = \sum_j \frac{C^2_F C^2_V}{4\pi^3} (g^2_{Vj} + 3g^2_{A_j}) m_B m_B (kT)^2 \mu e \left[ \pi^2 + \frac{E^2_\nu}{(kT)^2} \right] \frac{1}{1 + e^{-E_\nu/kT}}.
\]

We note that for nucleons-only matter, neutrino absorption on single nucleons can proceed only if the proton concentration exceeds some critical value in the range \((11 - 15)\%\) (see Ref. [27]). For matter with lower proton concentrations, neutrino absorption occurs on two nucleons. However, as shown in Ref. [10], neutrinos may be absorbed on single hyperons as long as the concentration of \( \Lambda \) hyperons exceeds a critical value that is less than 3% and is typically about 1%.

4. Results and discussions

The calculations above give the mean free path in a mixture in which all baryons are under degenerate conditions. Several effects of strong interactions must be included before the results in Eq. (6) through Eq. (9) may be utilized. The renormalization of the density of states at the Fermi surfaces results in the baryon masses \( m_B \) being replaced by \( m^*_B = p_F/v_F \). As pointed out by Iwamoto and Pethick [4], baryon-baryon interactions introduce further Fermi liquid corrections. Specifically, the axial vector interaction is significantly suppressed, which increases the mean free path of neutrinos in dense matter. So far, such an analysis has been restricted to pure neutron matter only. In a multicomponent system, this formalism must be extended to include correlations between the different species present. The effect of density correlations in the long wavelength limit can be related to the compressibility of the system [11].

In addition, the medium dependence of \( g_A \) must also be considered. For example, it has been argued [28] that \( |g_A| \) is quenched for nucleons in a medium. Whether \( g_A \) for hyperons is similarly affected is not known and is worth studying. Finally, depending on the momentum transfers involved, final-state interactions may modify the weak-interaction matrix element. Of the many corrections mentioned above, those due to the effective mass are the easiest to incorporate, since it would be contained even in a mean field description of the equation of state. Pending a more complete analysis, we consider below the modifications introduced by the multicomponent nature of the system incorporating only the effective mass corrections.

When neutrinos are degenerate, the relative abundances of the individual components do not play a significant role in determining the total mean free path. However, the extent to which the neutrino mean free path is altered by the presence of hyperons depends on the number of hyperonic species present. This may be illustrated by using the results for \( \mathcal{R}_{nc}^{(1)} = [C^2_V + 2C^2_A]_{\nu+B+\nu+B}/[C^2_V + 2C^2_A]_{\nu+n+\nu+n} \) and \( \mathcal{R}_{cc} = [C^2(g^2_V + 3g^2_A)]_{\nu+B_1+\nu+B_2}/[C^2(g^2_V + 3g^2_A)]_{\nu+n+\nu+p} \) listed in Tables 1 and 2. For example, in matter containing \( \Lambda, \Sigma^- \) and \( \Sigma^0 \) hyperons, the scattering mean free path...
is reduced by about 30 – 50% from its value in nucleons only matter. This reduction is achieved mainly by scattering on the Σ⁻ hyperon, since the Λ and Σ⁰ hyperons do not contribute to scattering in lowest order. Similarly, up to 50% reduction may be expected from absorption reactions. Again, reactions involving the Σ⁻ hyperon, particularly the one leading to the Σ⁰ hyperon, give the largest contribution.

The relative importance of the scattering and absorption reactions by degenerate neutrinos may be inferred by noting that

\[
\frac{\lambda_a}{\lambda_s} = \frac{1}{3 \mu_e} \left( \frac{\sum_i (C^2_{V_i} + 2C^2_{A_i})m^*_B}{\sum_j C^2(g^2_{V_j} + 3g^2_{A_j})m^*_B} \right).
\]

(10)

For a fixed lepton fraction and for neutrinos of energy \(E_\nu \sim \mu_\nu\), the factor \(\mu_\nu/\mu_e \leq 1\) (see Figs. 1 and 2). Inasmuch as the baryon effective masses are all similar in magnitude, the factor in the parenthesis above may be approximated by the ratio of the coupling constants. From the results in Tables 1 and 2, it is easy to verify that the factor containing the coupling constants is of order unity (but generally less than unity) and is not very sensitive to the number of hyperonic species present. We thus arrive at the result that the absorption reactions dominate over the scattering reactions by a factor of about three to five, even in the presence of hyperons. This conclusion is not affected by the inclusion of the baryon effective masses in the calculation of \(\lambda_a/\lambda_s\).

For non-degenerate neutrinos, and when \(k_B T \ll E_\nu v_{Fi}/c\),

\[
\frac{\lambda_a}{\lambda_s} = \frac{4}{15} \left( \frac{\sum_i (C^2_{V_i} + 4C^2_{A_i})p^2_{Fi}E^3_\nu(1 + e^{-E_\nu/kT})}{\sum_j C^2(g^2_{V_j} + 3g^2_{A_j})m^*_B m^*_B (kT)^2 \mu_e \left( \frac{E^2_\nu}{(kT)^2} \right)} \right).
\]

(11)

In this case, the abundances of the various particles play a significant role in determining whether or not the absorption reactions dominate over the scattering reactions. Some insight may be gained by examining the concentrations in Figs. 1 and 2, and also the ratio \(R^{(2)}_{nc} = [C^2_V + 4C^2_A]_{\nu+B \rightarrow \nu+B}/[C^2_V + 4C^2_A]_{\nu+n \rightarrow \nu+n}\) listed in Table 2.

5. Outlook

Neutrino signals in terrestrial detectors offer a means to determine the composition and the equation of state of dense matter. Many calculations of the composition of dense matter have indicated that strangeness-rich matter should be present in the core of neutron stars. Possible candidates for strangeness include hyperons, a \(K^-\) condensate, and quark matter containing \(s\)-quarks. Neutrino opacities for strange quark matter have been calculated previously. In this work, we have identified the relevant neutrino-hyperon scattering and absorption reactions that are important new sources of opacity. Much work remains to be done, however. Many-body effects which may introduce additional correlations in a mixture are worth studying. Full
simulations, of the type carried out in Refs. [8] including possible compositions and the appropriate neutrino opacities, will be taken up separately.

6. Acknowledgements

This work was supported by the U.S. Department of Energy under contract number DOE/DE-FG02-88ER-40388. We thank Jim Lattimer for helpful discussions.

7. References

[1] A. Burrows and J. M. Lattimer, Astrophys. Jl., 307 (1986) 178; A. Burrows, Ann. Rev. Nucl. Sci., 40 (1990) 181; and references therein.
[2] K. Hirata, et. al., Phys. Rev. Lett., 58 (1987) 1490; R. M. Bionta, et. al., Phys. Rev. Lett., 58 (1987) 1494.
[3] K. Sato, Prog. Theor. Phys. 53 (1975) 595.
[4] T. J. Mazurek, Astrophys. Space Sci. 35 (1975) 117.
[5] M. Prakash, et. al., Phys. Rep (1995) to be published; and references therein.
[6] V. Thorsson, M. Prakash and J. M. Lattimer, Nucl. Phys. A572 (1994) 693.
[7] M. Prakash, J. Cooke and J. M. Lattimer, SUNY preprint, SUNY-NTG-95-1.
[8] W. Keil and H. T. Janka, Astronomy and Astrophysics, 1994, to be published.
[9] O. V. Maxwell, Astrophys. Jl. 316 (1987) 691.
[10] M. Prakash, et. al., Astrophys. Jl. 390 (1992) L80.
[11] R. F. Sawyer, Phys. Rev. D11 (1975) 2740; Phys. Rev. C40 (1989) 865.
[12] D. L. Tubbs and D. N. Schramm, Astrophys. Jl. 201 (1975) 467.
[13] D. Q. Lamb and C. J. Pethick, Astrophys. Jl. 209 (1976) L77.
[14] D. Q. Lamb, Phys. Rev. Lett 41 (1978) 1623.
[15] D. L. Tubbs, Astrophys. Jl. Suppl. 37 (1978) 287.
[16] S. Bludman and K. Van Riper, Astrophys. Jl. 224 (1978) 631.
[17] R. F. Sawyer and A. Soni, Astrophys. Jl. 230 (1979) 859.
[18] N. Iwamoto, Ann. Phys. 141 (1982) 1.
[19] N. Iwamoto and C. J. Pethick, Phys. Rev. D25 (1982) 313.
[20] B. T. Goodwin and C. J. Pethick, Astrophys. Jl. 253 (1982) 816.
[21] A. Burrows and T. J. Mazurek, Astrophys. Jl. 259 (1982) 330.
[22] B. T. Goodwin, Astrophys. Jl. 261 (1982) 321.
[23] S. W. Bruenn, Astrophys. Jl. Suppl. 58 (1985) 771.
[24] L. Van Den Horn and J. Cooperstein, Astrophys. Jl. 300 (1986) 142.
[25] J. Cooperstein, Phys. Rep. 163 (1988) 95.
[26] A. Burrows, Astrophys. Jl. 334 (1988) 891.
[27] J. M. Lattimer, et. al., Phys. Rev. Lett. 66 (1991) 2701.
[28] D. H. Wilkinson, Phys. Rev. C7 (1973) 930; M. Rho, Nucl. Phys. A231 (1974) 493; G. E. Brown and M. Rho, Phys. Rev. Lett. 66 (1991) 2720.
[29] G. E. Brown and H. A. Bethe, Astrophys. Jl. 423 (1994) 659.
This figure "fig1-1.png" is available in "png" format from:

http://arxiv.org/ps/nucl-th/9508009v1
This figure "fig1-2.png" is available in "png" format from:

http://arxiv.org/ps/nucl-th/9508009v1