

Many-body corrections to charged-current neutrino absorption rates in nuclear matter

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Including nucleon-nucleon correlations due to both Fermi statistics and nuclear forces, we have developed a general formalism for calculating the charged-current neutrino-nucleon absorption rates in nuclear matter. We find that at one-half the nuclear density many-body effects alone suppress the rates by a factor of 2 and that the suppression factors increase to ~ 5 at $4 \times 10^{14} \text{ g cm}^{-3}$. The associated increase in the neutrino-matter mean free paths parallels what we have found for neutral-current interactions and may open up interesting possibilities in the context of the delayed supernova mechanism and protoneutron star cooling. [S0556-2813(99)05601-0]

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I. INTRODUCTION

The neutrino absorption and scattering opacities in the post-shock core of a supernova, in which nuclei are largely disintegrated into nucleons, determine the duration, spectrum, and flavor distribution of the emerging neutrino pulse. The neutrino-matter interaction rates can be related to the space- and time-dependent correlations among the set of density operators for the separate nuclear constituents (to find the Gamow-Teller parts we must consider separate spin-up and spin-down densities). In the case of neutral-current interactions [1], there is an instructive limit, which also provides an estimate of the effects, in which the combined limits of large nucleon mass and small neutrino energy allow the use of long-wavelength limits of equal-time correlation functions, in turn expressible in terms of the second derivatives of an energy density functional with respect to various densities. This approach is the direct multichannel generalization of the familiar results for light scattering from the thermal density fluctuations in a fluid, where it is the compressibility that determines the long-wavelength opacity, and it was used in Refs. [2] and [3] to find significant reductions of the neutral-current opacity in certain regions. In Burrows and Sawyer [1], an approach based on ring graphs was used to encompass these results and to extend them to domains in which the equal-time and long-wavelength limits are not clearly applicable.

The use of the equal-time and long-wavelength limits to express correlation functions in terms of static susceptibilities cannot be extended to the charged-current interactions when there is a large chemical potential difference between protons and neutrons. In the present work, we give a theoretical framework for addressing these opacities, based on summing ring graphs, together with the results of calculations with input parameters taken from the current phenom-

enology of nuclear matter. A complementary development can be found in Reddy *et al.* [4].

II. MANY-BODY FORMALISM FOR CHARGED-CURRENT RATES

The charged-current interactions of the nuclear medium are determined by the statistical averages involving the nucleonic charged-current operator $j_\mu = \bar{\psi}_p \gamma_\mu (1 - g_A \gamma_5) \psi_n$ and its Hermitian conjugate

$$W_{\mu\nu}(q, \omega) = -iZ^{-1} \int d^4x e^{-i\mathbf{q}\cdot\mathbf{x}} e^{i(\omega + \hat{\mu})t} \times \text{Tr} \{ e^{-\beta(H - \sum \mu_i N_i)} [j_\mu(\mathbf{x}, t), j_\nu^\dagger(0, 0)] \} \theta(t), \quad (1)$$

where Z is the partition function and $\hat{\mu}$ is the chemical potential difference, $\hat{\mu} = \mu_n - \mu_p$. The vector (\mathbf{q}, ω) is the momentum-energy transfer from the leptons to the medium. The transition rate, differential in angle and ω , for the reaction $\nu_e + n \rightarrow e^- + p$ is given in terms of $W_{\mu\nu}$ as

$$\frac{d^2\Gamma(\nu_e \rightarrow e^-)}{d\omega d\cos(\theta)} = (2\pi^2)^{-1} G_W^2 \frac{(E_1 - \omega)^2 [1 - f_e(E_1 - \omega)]}{1 - \exp[-\beta(\omega + \hat{\mu})]} \times \Lambda^{\mu\nu}(q, \omega) \text{Im} W_{\mu\nu}(q, \omega), \quad (2)$$

where

$$\Lambda^{\mu\nu} = (4E_1 E_2)^{-1} \text{Tr} [\hat{p}_1 (1 - \gamma_5) \gamma^\mu \hat{p}_2 \gamma^\nu (1 - \gamma_5)], \quad (3)$$

$q = [E_1^2 + (E_1 - \omega)^2 - 2E_1(E_1 - \omega)\cos\theta]^{1/2}$, and E_1 is the incident neutrino energy. To verify the factors in Eq. (2) replace the commutator in Eq. (1) by the unordered product, the $\theta(t)$ function by 1/2 in order to capture the imaginary part, and the statistical factor in the denominator by unity. Then, the introduction of a complete set of states between the current operators gives the inclusive $\nu_e \rightarrow e^-$ rate. Recalling that the Heisenberg picture for the density operators is de-

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fined with respect to the operator $H - \sum \mu_i N_i$, we note that the $e^{i\hat{\mu}t}$ factor in Eq. (1) will be canceled before the time integral is performed. The other term in the commutator gives the same answer, but with an additional factor $-\exp[-\beta(\omega + \hat{\mu})]$. The statistical factor in the denominator of Eq. (2) cancels when the two terms are combined.

Taking the nucleons as nonrelativistic, which is allowable up to about twice the nuclear density, we note that the components W_{i0} and W_{0i} vanish, and that the space tensor W_{ij} derives from the axial-vector current and can be written in the form

$$W_{ij}(\omega, \cos \theta) = g_A^2 W_A(\omega, \cos \theta) \delta_{ij} + g_A^2 W_T(\omega, \cos \theta) q_i q_j. \quad (4)$$

We define $W_{00} = W_V$, as it comes entirely from the vector current. Combining with Eq. (3) and calculating the lepton trace, we now have

$$W_{\mu\nu} \Lambda^{\mu\nu} = 2(1 + \cos \theta) W_V + 2(3 - \cos \theta) g_A^2 W_A + 2\omega^2(1 + \cos \theta) g_A^2 W_T. \quad (5)$$

We calculate the three functions W_V , W_A , and W_T in the ring approximation. For these graphs, a unit of charge is passed in each nuclear interaction in the expansion of the charged-current correlation function of Eq. (1). If the potential between nucleons a and b conserves isotopic spin, this means that only the terms containing the isotopic operator $\vec{\tau}^{(a)} \cdot \vec{\tau}^{(b)}$ enter. We assume a potential in momentum space,

$$V^{a,b}(q, \omega) = \vec{\tau}^a \cdot \vec{\tau}^b [v_1(q) + \vec{\sigma}^a \cdot \vec{\sigma}^b v_2(q) + \vec{\sigma}^a \cdot \vec{q} \vec{\sigma}^b \cdot \vec{q} v_3(q, \omega)]. \quad (6)$$

We have inserted in the v_3 term a dependence on energy transfer, ω , in order to accommodate the one-pion exchange force,

$$v_3(q, \omega) = -f^2 m_\pi^{-2} (q^2 + m_\pi^2 - \omega^2)^{-1}, \quad (7)$$

with $f^2 \approx 1$.

For the other two potentials, we follow the development of [1] in taking zero-range forms fitted to the Landau parameters of Fermi liquid theory and Ref. [5] and obtain $v_2 = 3.4 m_n / m^* \times 10^{-5} \text{ (MeV)}^{-2}$ and $v_1 = 1.88 m_n / m^* \times 10^{-5} \text{ (MeV)}^{-2}$, where m^* is the effective nucleon mass in the medium at nuclear density. We are assuming isospin-invariant forms, despite the fact that our application will be to unsymmetric matter. This appears to us to be the state of the art. It is not totally satisfactory, in view of the fact that the v_1 and v_2 terms are phenomenological forms that obtain in nuclear matter rather than forms based on elementary meson exchanges. Note that in the case of nonsymmetric matter isospin symmetry is still broken through the polarization functions. We define the nucleon-charge-raising polarization part $\Pi(q, \omega)$, in parallel with Eq. (1), in terms of the retarded commutator of the density $n_c = \psi_p^\dagger \psi_n$ and its Hermitian conjugate,

$$\begin{aligned} \Pi(q, \omega) &= -iZ^{-1} \int d^4x e^{-i\mathbf{q} \cdot \mathbf{x}} e^{i(\omega + \hat{\mu})t} \\ &\times \text{Tr}[e^{-\beta(H - \sum \mu_i N_i)} \{n_c(\mathbf{x}, t), n_c^\dagger(0, 0)\}] \theta(t). \end{aligned} \quad (8)$$

Note that the function that gives the vector part of the rate, $W_V(q, \omega)$, is exactly given by Π . The ring approximation [6] for Π is now

$$W_V(q, \omega) = \Pi(q, \omega) = \frac{\Pi^{(0)}(q, \omega)}{1 - 2v_1(q)\Pi^{(0)}(q, \omega)}, \quad (9)$$

where $\Pi^{(0)}$ is the polarization function in the absence of interactions. The factor of 2 multiplying the potential comes from the isospin operator in the potential (6). The spin-dependent parts of the potential do not contribute.

For the axial contribution, the operative nonrelativistic forms of the current are the operators $n_c^i = \psi_p^\dagger \sigma^i \psi_n$, and the analog to Eq. (8) will be a tensor with indices i, j . In the absence of interactions, this tensor is given by $\delta_{i,j} \Pi^{(0)}$. Thus, summing the axial chain to get $W_A(q, \omega)$ we obtain exactly the same structure as Eq. (8) with v_2 replacing v_1 , but with the same function $\Pi^{(0)}$,

$$W_A(q, \omega) = \frac{\Pi^{(0)}(q, \omega)}{1 - 2v_2(q)\Pi^{(0)}(q, \omega)}. \quad (10)$$

The potential v_3 does not enter Eq. (10), even though it couples to the axial-vector terms, since any ring graph chain in which at least one v_3 participates becomes a contribution to W_T . Elementary combinatorics for the tensor chain gives

$$W_T(q, \omega) = \frac{[\Pi^{(0)}(q, \omega)]^2 v_3(q)}{1 - 2[v_2(q) + q^2 v_3(q, \omega)]\Pi^{(0)}(q, \omega)}. \quad (11)$$

The requisite polarization function is given by

$$\Pi^{(0)}(q, \omega) = -2 \int \frac{d^3p}{(2\pi)^3} \frac{f(|\mathbf{p}|, \mu_n) - f(|\mathbf{p} + \mathbf{q}|, \mu_p)}{\omega + \epsilon_{\mathbf{p}} - \epsilon_{\mathbf{p} + \mathbf{q}} + i\eta}, \quad (12)$$

where the functions f are the nuclear Fermi occupation functions for the indicated momenta and chemical potentials. We conceptually extend Eq. (12) to include the average potentials that the nucleons experience in the medium, which we define as $V_{p,n}$, by making the replacements $(\mathbf{p})^2/(2m) \rightarrow (\mathbf{p})^2/(2m) + V_n$ and $(\mathbf{p} + \mathbf{q})^2/(2m) \rightarrow (\mathbf{p} + \mathbf{q})^2/(2m) + V_p$, both in the denominator and in the distribution functions in the numerator. However, we suppose that we are starting with a table of densities and temperatures (T) from an equation of state that already takes into account the potentials $V_{p,n}$. If we utilize Fermi distributions in which the chemical potentials are derived from the input densities using the free-particle relations, these average potential corrections are automatically included and the parameters $V_{p,n}$ do not appear explicitly in the formalism. To do the computations, we use the following form for $\Pi^{(0)}$:

$$\Pi^{(0)}(q, \omega) = \frac{m^2}{2\pi^2 q \beta} \left(\int_{-\infty-i\epsilon}^{\infty-i\epsilon} ds s^{-1} \ln[1 + e^{-(s+Q_+)^2 + \beta\mu_n}] + \int_{-\infty+i\epsilon}^{\infty+i\epsilon} ds s^{-1} \ln[1 + e^{-(s+Q_-)^2 + \beta\mu_p}] \right), \quad (13)$$

where

$$Q_{\pm} = \left(\frac{m\beta}{2} \right)^{1/2} \left(\mp \frac{\omega}{q} + \frac{q}{2m} \right), \quad (14)$$

which gives

$$\text{Im } \Pi^{(0)}(q, \omega) = \frac{m^2}{2\pi\beta q} \ln \left[\frac{1 + e^{-Q_+^2 + \beta\mu_n}}{1 + e^{-Q_+^2 + \beta\mu_p - \beta\omega}} \right] \quad (15)$$

and

$$\text{Re } \Pi^{(0)}(q, \omega) = \frac{m^2}{2\pi^2 q \beta} \int_0^{\infty} \frac{ds}{s} \ln \left[\frac{1 + e^{-(s+Q_+)^2 + \beta\mu_n}}{1 + e^{-(s-Q_+)^2 + \beta\mu_n}} \right] + (\omega \rightarrow -\omega, \mu_n \rightarrow \mu_p). \quad (16)$$

The imaginary part of the polarization (15) can be obtained by direct integration and is the same as that given in Refs. [1,7]. Given this, the full function (13) can be verified by checking the analytic properties in the ω plane. We re-emphasize that in Eqs. (15) and (16) μ_n and μ_p are derived from given particle densities via the free-particle prescription, and that despite this, these equations fully, if implicitly, incorporate the interaction potentials. However, Reddy, Prakash, and Lattimer [7] and Burrows and Sawyer [1] note a substantial effective mass correction to the rates. Obviously, a change in the mass parameter is not completely subsumed in the chemical potential-density relation and we still have to take into account its effects in the kinematics and on the particle density-chemical potential relation itself. Nevertheless, given an effective mass, our formalism easily allows it to be incorporated merely by substituting it for the nucleon mass in the normal way, wherever it explicitly appears.

We look at the results of medium interactions for two sets of conditions typical of the dense interior of a protoneutron star, but at different times: first for an early time ($t < 5$ s), when the lepton number is large, and second at a later time ($t > 10$ s), after which the lepton excess has largely radiated away [8].

III. LEPTON-RICH ERA

The densities of the various species are such that at zero temperature the inequality $2p_F^{(p)} > |p_F^{(n)} - p_F^{(v)}|$ holds. This inequality allows the single-nucleon process to proceed at full strength. That is to say, at low temperatures the function $\text{Im } \Pi^{(0)}(q, \omega)$, Eq. (15), is large in the region of (q, ω) defined by the two conditions: (1) the neutrino energy E_1 is near the neutrino Fermi surface; (2) the electron energy $E_1 - \omega$ is near the electron Fermi surface. These conditions come from the occupation factors that enter the expression for the total rate.

TABLE I. The total suppression factors ($S_{A,V}$) for the process $\nu_e + n \rightarrow e^- + p$, for a profile in an early post-bounce model generated by Burrows, Hayes, and Fryxell [9]. The suppression factors for the vector and the axial-vector terms are shown separately. These suppression factors are derived by multiplying the rate by the neutrino occupation function and integrating over neutrino energies. A nucleon effective mass of $0.75m_n$ is assumed.

ρ (g cm $^{-3}$)	Y_ν	T (MeV)	Y_e	S_A	S_V
3.94×10^{14}	0.077	5	0.289	0.140	0.269
3.68×10^{14}	0.078	5	0.294	0.144	0.275
3.08×10^{14}	0.077	5	0.297	0.157	0.291
1.65×10^{14}	0.064	10	0.275	0.228	0.381
2.66×10^{13}	0.01	15	0.282	0.670	0.775
1.40×10^{13}	0.067	15	0.258	0.790	0.840

In this region, we use Eqs. (9), (10), and (11) to calculate the modification factors in the medium. We have estimated the contribution of the tensor term (11) relative to the two other terms and conclude that it is less than 10% of the total in the cores of protoneutron stars and supernovas. For the vector and axial-vector terms, we compute suppression factors S_A and S_V defined as the ratio of the rates calculated with the nuclear interactions to those calculated without. This is done by substituting Eqs. (9) and (10) into Eq. (5), multiplying Eq. (2) by the neutrino occupation function, and integrating over neutrino energies. In Table I, we give results using a post-bounce supernova profile taken from Burrows, Hayes, and Fryxell [9]. As seen in the table, the Gamow-Teller suppression factors are larger than the Fermi suppression factors. Furthermore, the degree of many-body suppression increases with density and decreases with temperature. Importantly, the magnitude of the effect above 10^{14} g cm $^{-3}$ is large, ranging from a factor of 2 at 10^{14} g cm $^{-3}$ to a factor of ~ 5 near 4×10^{14} g cm $^{-3}$. Correspondingly, the neutrino-matter absorption cross sections decrease with density. Since it has recently been shown [1] that the neutral-current scattering rates at high density are also reduced, we conclude that post-bounce supernova cores are significantly more transparent than previously believed. As demonstrated in [1], this enhanced transparency can translate at intermediate times (> 500 – 2500 ms after bounce) into higher neutrino luminosities, which thereby may be more efficient at reenergizing or powering a stalled supernova shock. It is not only suppression at supranuclear densities that is germane to the luminosities after 1 to a few seconds. Equally important is the fact that such suppressions appear also near and below nuclear densities (see Table I).

Integrating Eq. (2) over $\cos \theta$, the distribution of the energy transfer, ω , to and from the nucleons due to the process $\nu_e + n \rightarrow e^- + p$ in the lepton-rich era can be derived and is depicted in Fig. 1 for a variety of densities, from 10^{13} to 10^{15} g cm $^{-3}$. For these curves, the temperature is 5 MeV, the electron fraction Y_e is 0.26, the incident electron neutrinos are on their Fermi surfaces, and beta equilibrium is assumed. The highest curve on the right (that for $\rho = 10^{13}$ g cm $^{-3}$) ignores many-body effects, though it incorporates the full kinematics, and is included for comparison. As expected, the peak of the energy transfer is generally near $-\hat{\mu}$ (given in the figure caption), since the electron

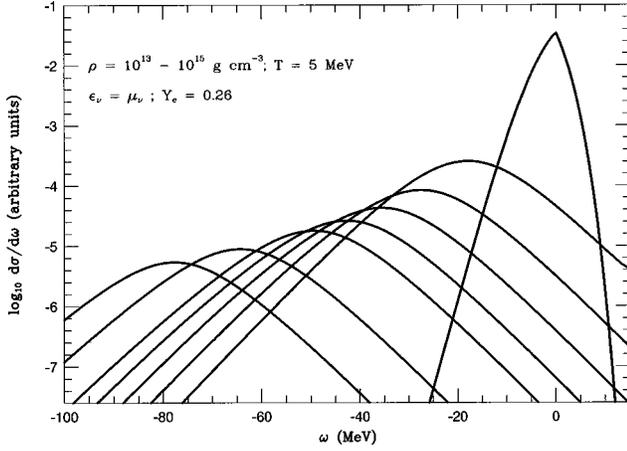


FIG. 1. The singly differential cross section in arbitrary units versus the energy transfer ω to the nucleons due to the process $\nu_e + n \rightarrow e^- + p$ in the lepton-rich era. ω is in MeV. We have integrated Eq. (2) over $\cos \theta$. The curves are for mass densities of 10^{13} , 10^{14} , 2×10^{14} , 3×10^{14} , 4×10^{14} , 5×10^{14} , 7.5×10^{14} , and 10^{15} g cm^{-3} . The temperature is 5 MeV, the electron fraction is 0.26, the incident electron neutrinos are on their Fermi surfaces, and beta equilibrium has been assumed. The 10^{13} g cm^{-3} curve does not include many-body effects, but does incorporate the full kinematics. For the other curves, $v_2 = 4.5 \times 10^{-5} (\text{MeV})^{-2}$, $v_1 = 1.76 \times 10^{-5} (\text{MeV})^{-2}$, and an effective mass at nuclear density of $0.75m_n$ was assumed. $\hat{\mu}$ is equal to 5.3, 17.4, 26.8, 34.8, 42.1, 48.7, 63.7, and 77.1 MeV for the densities depicted.

blocking factor in Eq. (2) puts the electrons on the electron Fermi surface and beta equilibrium requires that $\mu_e = \hat{\mu} + \mu_{\nu_e}$. (Note that $\hat{\mu}$ increases with density.) There is a modest spread in ω around the peak with approximately a Gaussian distribution. The width of this distribution scales with the temperature. Figure 1 demonstrates that Table I also reflects that the total cross sections, the integrals under the unnormalized curves, are decreasing functions of density.

IV. LEPTON-POOR ERA

As discussed in [10] and [11], as the trapped electron lepton number decays, we reach a configuration in which the neutrino absorption process discussed above dies almost completely for low temperatures. At the end of deleptonization, we have $p_F^{(\nu)} = 0$ [8]. The neutrinos then have thermal energies and we find that the proton fraction has decreased to the point that $2p_F^{(p)} \ll p_F^{(n)}$. In this case we cannot conserve momentum for the three degenerate species e^- , p , and n when the momentum of the neutrino is small and when we stay near the Fermi surface for the three other species. At low temperatures, the function $\text{Im} \Pi^{(0)}(q, \omega)$ of Eq. (15) is now exponentially small in the region of (q, ω) defined by the leptonic occupancies. Thus, the ring graphs, as defined in the previous section, give a negligible rate for the charged-current processes at low lepton number and temperature. The region of temperature and density in which this inhibition prevails depends on the interactions in the medium and, indeed, on the whole equation of state.

The mechanisms that have been proposed to estimate the rates in this domain depend on a spectator nucleon to transfer

the necessary momentum, either through a potential or through an assumed correlation [10–12]. Translated to graphs, these mechanisms involve the estimation of proper graphs for the polarization parts, where a proper graph is defined as a graph that cannot be cut into two disjoint parts by severing a single potential line. However, there are ‘‘ring corrections’’ to such graphs, in which the initial or final current vertices attach to a ring chain that then attaches to the proper polarization graphs. These ring corrections then have a big suppressive effect on the primary mechanism for momentum transfer, as one can see from the following argument.

We consider a proper polarization graph in which a nucleon-nucleon interaction has intervened to allow momentum conservation near the Fermi surfaces. That is to say, by Eq. (10), we have a term for the vector contribution W_V to the rate formula (5) that has a substantial imaginary part in the kinematically allowed region. We call this term $W_V^{(1)}$. We now take the sum of this term and the lowest-order term, $\Pi^{(0)}$, as the proper polarization part from which to construct the ring sum, obtaining

$$\begin{aligned} \text{Im } W_V(q, \omega) \\ \approx \frac{\text{Im } W_V^{(1)}(q, \omega)}{|1 - 2v_1[\Pi^{(0)}(q, \omega) + \text{Re } W_V^{(1)} + i \text{Im } W_V^{(1)}(q, \omega)]|^2}. \end{aligned} \quad (17)$$

In the estimates that follow, we omit the real part of $W_V^{(1)}$ from the denominator; it is easy to verify from Eq. (16) that, in contrast to the lowest-order imaginary part, the real part is not suppressed in the region of (q, ω) that is important in the reaction.

Similar considerations hold for the axial-vector part. In the lepton-poor era, we can drop the second term on the right-hand side of Eq. (5), since ω will be of the order of T , rather than of order 100–200 MeV, as it can be for electron neutrinos in the trapped neutrino era [13]. Then, we define a proper W_A and a contribution $W_A^{(1)}$ as in the above and write

$$\text{Im } W_A(q, \omega) \approx \frac{\text{Im } W_A^{(1)}(q, \omega)}{|1 - 2v_2[\Pi^{(0)}(q, \omega) + i \text{Im } W_A^{(1)}(q, \omega)]|^2}. \quad (18)$$

TABLE II. The total axial and vector suppression factors ($S_{A,V}$) for the transformation $\nu_e \rightarrow e^-$ during the lepton-poor era, for a density of 2.5×10^{14} g cm^{-3} and a neutrino chemical potential of zero. These suppression factors are derived by multiplying the rate by the neutrino occupation function and integrating over neutrino energies. A nucleon effective mass of $0.75m_n$ is assumed. See the text for details.

T (MeV)	Y_e	S_A	S_V
3	0.012	0.34	0.53
5	0.013	0.20	0.38
7	0.016	0.19	0.38
9	0.018	0.20	0.38

For the factors $\text{Im } W_{V,A}^{(1)}$ we take the minimal form consistent with avoiding a singularity in Eq. (2) at $\omega = -\hat{\mu}$:

$$\text{Im } W_{V,A}^{(1)} = c_{V,A}^{(1)}(\omega + \hat{\mu}). \quad (19)$$

This form meets the requirement for detailed balance that $\text{Im } \Pi$ be odd under the replacements $\omega \rightarrow -\omega$, $\hat{\mu} \rightarrow -\hat{\mu}$. We take the parameters $c^{(1)}$ to be sufficiently small for the integrated suppression factor to be independent of the $[c^{(1)}]^2$ that enter through the imaginary parts in the denominator function. In Table II we give, for the case of matter at nuclear density, the separate suppression factors for the vector and axial-vector rates generated by the uncorrected terms (19), under deleptonized conditions, $\mu_\nu = 0$. We note that the reductions are substantial, and conclude that the ring corrections should be added to any model that is used for the neutrino opacities during this era.

V. CONCLUSIONS

We have developed a formalism for incorporating the effects of many-body correlations on the charged-current rates of neutrino-matter interactions. This formalism reveals that these rates are considerably suppressed in the densest regions of proton-neutron stars and supernova cores. Assuming that the nucleons are nonrelativistic, our formalism incorporates the full kinematics of the interaction, Pauli blocking by final-state nucleons (protons), and correlation due to nucleon-nucleon interactions.

We have employed the ring approximation and assumed the near validity of Fermi liquid theory. It would be desirable to include ladder diagrams and to perform the calculations in the context of a better numerical method for solution of the nuclear equation of state (EOS), since the solution of the

EOS is intimately related to the derivation of the scattering and absorption rates. However, those who perform detailed nuclear EOS calculations and address many-body correlations in nuclear matter do not as yet provide the requisite spin and density structure functions, even for the static case.

These results for charged currents, when combined with the results from Burrows and Sawyer [1] for neutral currents, suggest that energy and lepton number will leak from supernova cores at a rate that is higher than heretofore estimated. This implies that the neutrino luminosities during the epoch after bounce for which the inner core is the major energy source ($>0.5 - 1.5$ s) will be enhanced for seconds, perhaps by as much as 50% [1], after which time they would be diminished, in keeping with total energy conservation. The consequences of this increased transparency for the neutrino-driven supernova explosion mechanism [9] may be interesting, but have yet to be clarified. Note that the true magnitude of this effect can in principle be constrained by the neutrino signal from SN1987A [14], but by how much depends upon its nascent neutron star mass and the actual nuclear equation of state, both poorly known. Importantly, using the duration and number of events to obtain such constraints is hampered by small number statistics. Be that as it may, it is clear that many-body effects provide new possibilities that supernova modelers should enjoy incorporating into the next generation of explosion calculations.

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