Neutrino Heating in Supernovae

W. C. Haxton
Institute for Nuclear Theory, Department of Physics, University of Washington, Seattle, Washington 98195
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I argue that standard descriptions of stellar collapse omit the primary mechanism for dissipative neutrino reactions in nuclear matter, nuclear excitation by neutral-current scattering. The nuclear heating rate, due primarily to muon- and tauon-neutrino excitation of giant resonance states, is on the order of 90 MeV/nucleon sec at a radius of 100 km. I discuss possible effects of both neutral- and charged-current neutrino heating in models of stellar collapse.

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According to the standard models of supernovae,[1-4] approximately 99% of the energy of the collapse is radiated in neutrinos. In this Letter, I discuss the inelastic interactions by which these neutrinos can deposit their energy in nuclear matter outside the neutrinosphere. I find that the primary dissipative mechanism during the late-time (i.e., neutron-star cooling) phase of the collapse is one that is not considered in standard supernova calculations: nuclear excitation by neutral-current neutrino scattering, particularly muon- and tauon-neutrino excitation of giant resonance states. I discuss the nuclear-physics arguments that lead to this result, and speculate on the possible consequences for models of stellar collapses. I also discuss charged-current heating during the first few milliseconds following shock-wave formation.

There are two standard scenarios for type-II supernovae, the prompt-explosion model[4] of the Stony Brook–Brookhaven group and the delayed-explosion model[2,3] of Wilson and his collaborators at Livermore. The issue that separates these models is the dissipation of the shock wave's energy as it propagates through the star. In the Stony Brook–Brookhaven model a light progenitor ($\lesssim 15 M_\odot$) with a small iron core ($\lesssim 1.35 M_\odot$) produces a sufficiently energetic shock to overcome losses due to nuclear dissociation and neutrino-pair production. This model requires a soft nuclear equation of state to produce an energetic shock wave. In the Livermore model the shock wave weakens and stalls at a radius of a few times 10^7 cm, but then, in some cases, is revived about 0.5 sec later by the charged-current neutrino "reheating" of free nucleons left in the wake of the shock. In each model the energy carried by the shock wave is a few percent of the gravitational binding energy of the remnant neutron star, about 3×10^{53} ergs, while the remaining, much larger fraction of the energy is radiated in neutrinos. The assumptions[5,6] made in each model about neutrino interactions with matter are similar: Neutrinos lose energy by their charged-current reactions on free protons and neutrons (and some nuclei), by charged-current ν_e and $\bar\nu_e$ scattering off electrons, and by neutral-current scattering off electrons. Muon and tauon neutrinos participate only through the last reaction, which has a very weak cross section

$$[\sigma_{\nu_\mu}=\sigma_{\nu_\tau}=\tfrac{1}{6}\,\sigma_{\nu_e}=1.4\times10^{-45}E_\nu/(1\text{ MeV})\text{ cm}^2].$$

The models include neutral-current elastic scattering off free nucleons and nuclei (indeed this leads to neutrino trapping for densities $\rho\lesssim 3\times10^{11}$ g/cm^3), but this process is not dissipative.

To set a benchmark for later discussions of inelastic neutral-current reactions, I summarize the standard treatment[2] of neutrino heating by charged-current reactions. I assume that the neutrinos are emitted from the neutrinosphere in a Fermi-Dirac distribution and move radially outward (so that the finite extent of the neutrinosphere is neglected). The rate of energy deposition due to (ν_e,e) reactions in a volume element outside the neutrinosphere is

$$\frac{dE_\nu^n}{dt}=(223\text{ MeV/nucleon sec})\left(\frac{1}{R_7}\right)^2\left[Y_nL_{52}(\nu_e)\left(\frac{T_{\nu_e}}{5\text{ MeV}}\right)^2+Y_pL_{52}(\bar\nu_e)\left(\frac{T_{\bar\nu_e}}{5\text{ MeV}}\right)^2\right], \tag{1}$$

where R_7 is the distance from the neutron-star center in units of 10^7 cm, $L_{52}(\nu_e)$ is the ν_e luminosity in units of 10^{52} ergs/sec, Y_n and Y_p are the mass fractions in free neutrons and protons, and T_{ν_e} is the temperature of the Fermi-Dirac distribution multiplied by Boltzmann's constant. In the late-stage reheating model, the region behind the shock wave where net positive heating occurs is relatively narrow,[6] bounded on the outside by a layer of undissociated ^{4}He and on the inside by hot ($T_m\gtrsim2$ MeV) baryonic matter where neutrino losses due to electron capture exceed the heating from Eq. (1).

I now consider the heating of nuclear matter. As in the case of baryonic matter, electron neutrinos can deposit their energy by charged-current reactions (ν_e,e^-) and $(\bar{\nu}_e,e^+)$. However, for nuclei like ^{4}He, ^{12}C, and ^{16}O, the allowed transition strength is a small fraction of that for the free nucleon and the nuclear thresholds are large. Thus the nuclear cross sections are small. A principal point of this paper is that a second process, inelastic neutrino scattering, generally provides a much more effective heating mechanism. This comes about for several reasons:

(1) All neutrino flavors participate in this process.

(2) Once the collapse has entered the cooling stage, the temperatures of the muon and tauon neutrinos rise to about twice that of the electron neutrinos. The reactions of these more energetic species are favored by phase space and are not as strongly affected by nuclear thresholds.

(3) Muon and tauon neutrinos from the high-energy tails of the Fermi-Dirac distributions can strongly excite the first-forbidden transitions to giant resonances. The sum rules that govern these transitions lead to energy-weighted cross sections that are roughly proportional to the atomic number A.[7] These contributions produce substantial cross sections even for closed-shell nuclei like ^{4}He and ^{16}O, where almost no Gamow-Teller or Fermi strength exists.

The heating rate in matter composed of nuclei is, in analogy with Eq. (1),

$$\frac{dE_\nu^A}{dt}=(31.6\ \text{MeV/nucleon sec})\left(\frac{1}{R_7}\right)^2\left[L_{52}(\nu_e)\left(\frac{5\ \text{MeV}}{T_{\nu_e}}\right)\frac{A^{-1}\langle\sigma_{\nu_e}^+E_{\nu_e}+\sigma_{\nu_e}^0E_{\text{ex}}^A\rangle_{T_e}}{10^{-40}\ \text{cm}^2\ \text{MeV}}\right.$$
$$+L_{52}(\bar{\nu}_e)\left(\frac{5\ \text{MeV}}{T_{\bar{\nu}_e}}\right)\frac{A^{-1}\langle\sigma_{\bar{\nu}_e}^-E_{\bar{\nu}_e}+\sigma_{\bar{\nu}_e}^0E_{\text{ex}}^A\rangle_{T_{\bar{e}}}}{10^{-40}\ \text{cm}^2\ \text{MeV}}$$
$$\left.+L_{52}(\nu_\mu)\left(\frac{10\ \text{MeV}}{T_{\nu_\mu}}\right)\frac{A^{-1}\langle\sigma_{\nu_\mu}^0E_{\text{ex}}^A+\sigma_{\bar{\nu}_\mu}^0E_{\text{ex}}^A\rangle_{T_\mu}}{10^{-40}\ \text{cm}^2\ \text{MeV}}\right], \tag{2}$$

where $\sigma_\nu^\pm$ and σ_ν^0 denote the charged-current and neutral-current cross sections, E_ν is the neutrino energy, E_{ex}^A is the nuclear excitation energy, and $\langle\ \rangle_T$ denotes an average over a normalized neutrino spectrum (again assumed to be Fermi Dirac). The last term, though written for muon neutrinos, gives the sum of the contributions from both muon and tauon neutrinos ($L_{\nu_\mu}=L_{\nu_\tau}$, $T_{\nu_\mu}=T_{\nu_\tau}$).

I have estimated these cross sections for ^{4}He, ^{12}C, ^{16}O, and ^{56}Fe in a series of shell-model calculations. In each case, sum-rule considerations governed the choice of model space and Hamiltonian. For ^{4}He complete sets fo nonspurious $2\hbar\omega$ and $3\hbar\omega$ states were generated by diagonalization of the Sussex interaction.[8] In a harmonic oscillator single-particle basis, the summation over these states preserves the Gamow-Teller (GT) sum rule and the sum rules for the first-forbidden axial charge, $\mathbf{r}$, and $[\sigma\otimes r]_J$ operators. In ^{12}C the positive- and negative-parity states were described by a complete set of $0\hbar\omega$ and $1\hbar\omega$ wave functions derived from the Cohen-Kurath[9] and the Millener-Kurath[10] interactions. The positive-parity diagonalization for ^{16}O was performed for a complete set of $2\hbar\omega$ states, with the Kuo[11] and Brown-Wildenthal[12] interactions supplementing those mentioned earlier. Because a full diagonalization of a $3\hbar\omega$ space was impractical, transitions to negative-parity states were modeled as simple one-particle, one-hole ($1p$-$1h$) excitations of a closed core. That is, a cruder description of the ^{16}O ground state was adopted to avoid violating the sum rules for the first-forbidden operators. Finally, the ^{56}Fe ground state was described as a $2p$-$2h$ excitation of a ^{56}Ni core, with the particle states confined to the $2p_{3/2}$ shell and the holes to the $1f_{7/2}$ shell. Transitions were then summed to the complete set of positive-parity $2p$-$2h$ and $3p$-$3h$ fp-shell states required by the GT sum rule. Unfortunately, the analogous calculation for the negative-parity transitions is not feasible. Instead, I repeated the negative-parity calculations for each nucleus in the Goldhaber-Teller model, which provides a simple description of the fifteen-dimensional SU(4) supermultiplet of giant resonances.[7] The ratios of the ^{4}He, ^{12}C, and ^{16}O shell-model neutral-current cross sections to those of the collective-model average 1.38 (and range from 1.27 to 1.54). I then used the collective-model results for ^{56}Fe, scaled by this correction factor. The inclusive-neutrino-reaction cross sections were evaluated as in a previous work.[13]

Calculations were performed for both charge-current and neutral-current cross sections for temperatures between 1 and 14 MeV. The neutral-current results for 5 MeV $< T <$ 14 MeV are given to an accuracy of better than 5% by

$$A^{-1}\langle\sigma_\nu^0E_{\text{ex}}^A+\sigma_{\bar{\nu}}^0E_{\text{ex}}^A\rangle_T=\alpha\left(\frac{T-T_0}{10\ \text{MeV}}\right)^\beta,$$

where α, β, and T_0 are given in Table I. At $T=10$ MeV, a typical temperature for late-stage muon and tauon neutrinos, this quantity is 0.41, 0.61, 0.52, and 0.57 for ^{4}He, ^{12}C, ^{16}O, and ^{56}Fe, respectively, in units of

TABLE I. Parametrizations of the energy-weighted neutral-current sections per nucleon.

	α (10^{-40} MeV cm^2)	β	T_0 (MeV)
^{4}He	1.24	3.82	2.54
^{12}C	1.30	2.83	2.36
^{16}O	1.38	3.14	2.66
^{56}Fe	1.00	1.76	2.73

TABLE II. The cross sections $A^{-1}\langle\sigma_{\nu_e}^{+}E_{\nu_e}\rangle_T$ (left-hand columns) and $A^{-1}\langle\sigma_{\bar{\nu}_e}^{-}E_{\bar{\nu}_e}\rangle_T$ in units of 10^{-42} MeV cm^2.

T (MeV)	^{4}He		^{12}C		^{16}O		^{56}Fe	
3	0.020	0.016	0.19	0.22	0.068	0.083	4.57	0.89
4	0.30	0.22	1.46	1.29	0.72	0.65	17.2	2.79
5	1.84	1.20	6.12	4.50	3.78	2.74	44.5	6.48
6	7.08	4.21	18.3	11.7	13.0	8.13	92.2	12.6

10^{-40} MeV cm^2. In Table II the corresponding charged-current results are given for a few ν_e temperatures near the expected value of about 5 MeV.

Charged-current cross sections in supernova codes are modeled in terms of effective free nucleons. However, the ^{4}He and ^{16}O cross sections are almost entirely determined by giant resonance transitions. These have no analog in the free nucleon, and grow much more steeply as a function of T than allowed cross sections. Even ^{12}C, despite having appreciable GT strength, has a cross section dominated by giant resonance transitions for $T \gtrsim 5$ MeV.

Neutral-current inelastic scattering off nuclei is not considered in standard supernova models. The energy-weighted cross sections at $T \sim 10$ MeV are dominated by giant resonance transitions for ^{4}He, ^{12}C, and ^{16}O. For ^{56}Fe, because of the strong GT transitions, they account for 55% of the neutral-current section. Muon- and tauon-neutrino interactions provide 92%, 82%, 87%, and 57% of the net heating for ^{4}He, ^{12}C, ^{16}O, and ^{56}Fe, respectively, if we assume comparable neutrino luminosities for each flavor and $T_{\nu_\mu} = 2T_{\nu_e} = 10$ MeV. For a mean luminosity of 4×10^{52} ergs/sec, the total heating rates for these nuclei are (56, 95, 77, and 130)R_7^{-2} MeV/nucleon sec, values that are (6.3-14.6)% that of Eq. (1) when $Y_n + Y_p = Y_N = 1$.

I now speculate on the possible consequences of these results.

(i) In the early stages of collapse, neutral-current scattering will combine with other inelastic processes to enhance the downscattering (thermalization) of electron neutrinos, with the resulting Pauli blocking helping to maintain a larger core lepton fraction. Naively one would expect this to produce a somewhat larger homologous core, leading both to a stronger shock wave at the bounce and less overlying iron for the shock wave to photodisintegrate. Unfortunately, as Cooperstein points out,[14] downscattered neutrinos can escape more readily because of their longer mean free paths. Furthermore, the nuclear heating accompanying downscattering would liberate free protons that rapidly consume electrons. Both of these phenomena will increase the electron-capture rate. Thus it would seem that rather careful calculations are needed to assess the net effect of the enhanced downscattering.

(ii) A few milliseconds after core bounce the shock wave moves rapidly ($v \sim 0.1c$) through the neutrinosphere at $R_7 \sim 0.5$. At this point the shock wave's energy is being dissipated by the photodissociation of ^{56}Fe nuclei falling through the shock front, and by neutrino production. The sharp burst of relatively low-energy ν_e's associated with the shock's penetration of the neutrinosphere carries about 3×10^{51} ergs. As the duration of the pulse is about 3 msec, the peak luminosity ($\sim 5\times10^{53}$ ergs/sec) is extraordinary. I now argue that a significant fraction of the energy carried by the burst neutrinos is not lost, but instead goes into preheating iron nuclei outside the shock front; i.e., the iron outside the shock wave is not transparent to these neutrinos.

Consider the preheating of an infalling volume of ^{56}Fe that meets the shock wave near the neutrinosphere. I assume that this volume has fallen inward with a velocity of about 0.6 of free fall, achieving $v \sim 0.1c$ at $R_7 \sim 0.5$. The accelerating core mass is taken to be $1.0M_\odot$. For the electron capture burst I use the profile of Fig. 5(1) of Ref. 6 ($25M_\odot$ collapse, model C). I assume an average ν_e temperature of 4 MeV, a value consistent with some recent calculations.[14] The corresponding energy-weighted cross section in iron is

$$A^{-1}\langle\sigma_{\nu_e}^{+}E_{\nu_e} + \sigma_{\nu_e}^{0}E_{\rm ex}^{A}\rangle = 0.19\times10^{-40}\ \text{MeV cm}^2;$$

most of this is due to the allowed charged-current reactions, which are enhanced by the strong Coulomb field for $Z=26$. However, the first-forbidden corrections to this cross section and the neutral-current contribution are significant, each increasing the net heating by 12%. I ignore contributions from other flavors, which have much lower fluxes at this time.

With the peak of the ν_e burst fixed in time at shock breakout, the calculated preheating of matter just in front of the shock wave ($R_7 \sim 0.5$) is 1.46 MeV/nucleon. Most of the heating occurs on a millisecond time scale in a narrow region outside the shock front. Similarly, infalling matter meeting the shock wave at $R_7 = 0.3$, 0.7, and 1.0 is preheated to 2.27, 1.07, and 0.75 MeV/nucleon, respectively. The results are not sensitive to changes in the shock velocity (which was varied by a factor of 2 about $0.1c$), and only somewhat sensitive to changes in the timing of the ν_e burst of up to ±1 msec. There is sensitivity to the neutrino burst temperature: The use of

$T_{\nu_e}=3$ MeV multiplies the above results by 0.36, while $T_{\nu_e}=5$ MeV yields a scale factor of 2.04.

These estimates do not include electron and neutrino blocking factors, though they should be small outside the neutrinosphere. Also, as the ν_e burst begins, a temperature below 4 MeV might be more appropriate, as the lower-energy neutrinos are liberated first. This will reduce the heating of infalling material that crosses the shock front at early times (such as the $R_7=0.3$ results above). However, probably the most serious flaw is one that makes the heating estimates too small: As neutrinos interact outside the shock front, the opacity of this material increases. If we assume that all of the deposited energy goes into liberating neutrons and protons (at a cost of about 9 MeV/nucleon), the opacity doubles by the time the heating reaches 1 MeV/nucleon. For the conditions described above, free neutrons generated in the heating process enhance the total heating by factors of 2.3, 1.9, and 1.6 for material crossing the shock front at $R_7=0.5$, 0.7, and 1.0, respectively.

As preheating means that less energy is required from the shock wave to dissociate the infalling Fe, the shock wave retains more of its strength while propagating through the iron core, thus enhancing the prospects for an explosion. It is clearly important to determine whether the rough arguments above hold up in supernova code calculations.

(iii) In the late-stage model, neutrino heating will occur in the ^{4}He zone just behind the stalled shock wave and will be dominated by muon and tauon neutrinos. In Table I of Ref. 2, the mass points just behind the shock front at times $t=0.417$ (when the shock stalls), 0.433, and 0.464 sec have nucleon mass fractions of $Y_N=0.010$, 0.011, and 0.064, respectively. As nuclear heating dominates for $Y_N \lesssim 0.06$, the neutral-current process will accelerate the initial heating, then become less important as Y_N climbs above this value. Also, there is some preheating of material falling through the shock front but, unlike in (ii), this effect is small [$\sim$(2–3)% of the ^{4}He dissociation energy] because of the larger radius ($R_7 \sim 3$). Nuclear preheating by muon and tauon neutrinos could become more important if the stalled shock wave recedes to a radius much smaller than $R_7 \sim 3$.

(iv) Perhaps the most significant effect of the neutral-current reactions may be in the alteration of the chemical environment in the mantle of the star. With nuclear cross sections of a few times 10^{-41} cm^2, large numbers of protons, neutrons, and neutrino spallation products will be produced. Certain products, such as the fluorine produced by spallation in the neon shell, can provide important constraints on the flux and temperature of the muon and tauon neutrinos.[15] The contribution to r-process nucleosynthesis also appears to be interesting.[16]

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