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40561-072-INT92-00-23

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Talk presented at the Benjamin Franklin Symposium

in celebration of the Discovery of the Neutrino

Drexel University, Philadelphia, PA, April 29-May 1, 1992

PREPARED FOR THE U.S. DEPARTMENT OF ENERGY

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ABSTRACT

Neutrinos provide an important diagnostic of behavior deep inside the solar core. They also play a major role in the evolution of red giants, horizontal branch stars, and core-collapse supernovae. I discuss some of the interesting neutrino physics issues that have arisen in studies of stellar evolution.

1. Introduction

It is a great pleasure to participate in this symposium in honor of Fred Reines. Fred's lifelong fascination with the neutrino illustrates more clearly than I can express the deep connections between terrestrial experiments and astrophysical neutrinos. Fred's discovery of neutrinos in a reactor experiment[1] led to his appreciation of the importance, both in nuclear/particle physics and astrophysics, of measuring the flux and spectra of neutrinos from our sun[2] and from distant supernovae.[3] As this symposium illustrates, this quest remains an exciting one.

2. Solar and Red Giant Neutrinos

Previous talks have described the importance of solar neutrino research in testing the standard solar model predictions of conditions in the sun's core. Thanks to Mikheyev, Smirnov, and Wolfenstein[4], we also appreciate the sun as a marvelous regenerator, enhancing the effects of oscillations through adiabatic level crossings, and thereby greatly extending the ranges of neutrino masses and mixing angles that we can plausibly explore. Despite all of this, the impact of neutrino physics on our sun's evolution is not a dramatic one: the energy carried away by neutrinos represents about 2% of the core's output, while the remainder is transported to the solar surface by radiation and convection. Furthermore, the only solar neutrino flux we have definitely observed, the ^{8}B neutrinos, is of relatively recent origin, having been less than half its contemporary value when the sun was 80% of its current age, according to the standard solar model.[5] Thus, through most of the sun's lifetime, these neutrinos have played no role.

In contrast, neutrino emission can be the dominant mechanism for energy transport in other classes of stars, such as red giants (RG). A red giant is a star that has exhausted the hydrogen in its core, leaving a dense ^{4}He core that supports itself by the degeneracy pressure. Surrounding this core is a hydrogen mantle, with H burning taking place in a narrow shell just outside of the ^{4}He core. As the star evolves, more ^{4}He is added to the core until finally it reaches the density/temperature conditions that permit ^{4}He ignition. The mass of the core at ignition is crucial in determining the subsequent evolution of the star along the horizontal branch (HB).

Ignition depends delicately on conditions in the RG core, since the rate for $3\alpha \rightarrow {}^{12}$C varies roughly as $\sim \rho^2\, T^{30-40}$, where ρ and T are the density and temperature of the core. The primary mechanism for transporting energy out of the core is the decay of a plasmon into a $\nu\bar{\nu}$ pair.[6] This neutrino physics determines the point where the helium flash occurs, thereby fixing the core mass that governs the subsequent horizontal branch evolution, including the star's brightness and longevity on the HB.

Particle and nuclear astrophysicists have long understood that RG and HB age and brightness observations could be used to constrain anomalous cooling mechanisms that might compete with or enhance $\nu\bar{\nu}$ emission. One example is axion emission via Compton production off electrons, Primakoff production off nuclei, or emission by nuclear decays. Arguments based on RG energetics and on our understanding of the metallicity dependence of RG and HB evolution have been used to exclude axions with masses $\gtrsim 1$ eV.[7] Similarly, our most stringent limits on diagonal and transition neutrino magnetic and electric dipole moments comes from the direct plasmon $\rightarrow \nu\bar{\nu}$ cooling that such moments would generate. The recent analysis of Raffelt[8] yielded $\mu_{ij} \lesssim 3 \cdot 10^{-12} \mu_B$, where μ_B is an electron Bohr magneton. This in turn forces one to assume enormously strong magnetic fields in the sun's interior if one wishes to invoke neutrino magnetic moments as a possible source of variability in the solar neutrino flux.

3. The Supernova Mechanism

Another example where neutrinos play a major role in stellar evolution is core-collapse supernovae. Because of Fred's role in the first detection of supernova neutrinos[3], this seems like an appropriate focus for my talk. Neutrino physics is important to the evolution of the presupernova star, affecting the star's entropy budget and determining parameters, such as the mass and lepton number of the homologous core, that are crucial in theoretical models of the supernova explosion. After core bounce neutrinos play a major role in nucleosynthesis (because of nuclear reactions very similar to those recorded in the IMB and Kamioka II detectors) and very likely in powering the explosion itself.

Consider a massive star, $M \gtrsim 10 M_\odot$, burning its core hydrogen in hydrostatic equilibrium. If one takes $M \sim 25 M_\odot$ as typical of a supernova progenitor, this burning phase is completed in about 7 My, after which the core hydrogen is exhausted, having been burned to ^{4}He. The core, lacking an energy source to maintain the electron gas pressure, contracts until conditions are reached where ^{4}He ignition can occur. This cycle, fuel exhaustion, contraction, and subsequent ignition of the ashes of the previous burning cycle, repeats through carbon, neon, oxygen, and silicon burning, with this final burning stage producing in about one day an inert iron core.

Since iron is the most stable nucleus, no further burning reactions are possible.

Thus nothing is available to halt the contraction of the iron core ($\sim$ 1.2 - 1.5 $M_\odot$). As the collapse of the iron core proceeds, matter is heated, leading to the photodissociation of iron into alphas and nucleons. This in turn accelerates the weak capture of electrons by free protons, producing neutrons and escaping ν_es, and further reducing the gas pressure. Thus the collapse is rapid, proceeding at about 0.6 of the free-fall velocity according to numerical simulations.[9,10]

At a density $\rho \sim 10^{12}$ g/cm^3 the emitted ν_es, which up to this point have streamed through the star's mantle, carrying away some of the star's lepton number, become "trapped" by their neutral current scattering off nuclei. That is, elastic scattering leads to a diffusion time for the neutrinos that is longer than the time required for the collapse to be completed. Trapping thus means that the gravitational energy released by the remainder of the collapse remains in the core until well after the core bounce.

The velocity of sound in nuclear matter increases with increasing density. The homologous core is the inner portion of the iron core, usually 0.6 - 0.9 $M_\odot$, where the sound velocity exceeds the infall velocity. Since the sound velocity is then sufficient to smooth out pressure variations on the timescale of the collapse, it is not surprising that the density profile of the homologous core is unchanged during the collapse.

As nuclear density is reached ($\rho \sim 3 \cdot 10^{14}$ g/cm^3), the nuclear equation of state takes over, providing a new source of pressure that can halt the collapse. The innermost shell of matter at the center of the homologous core compresses to about twice nuclear density (depending on the details of the equation of state), with the rebound producing a pressure wave that starts to propagate back out through the homologous core. Successive shells repeat this process, producing new waves that follow the first. These collect at the edge of the homologous core, where $v_{\rm sound} \sim v_{\rm infall}$. A shock wave is released as the edge of the homologous core comes to rest, and begins to propagate through the outer portion of the iron core.

As the shock wave propagates outward, iron falling through the shock front is heated and melts, consuming about 8 MeV/nucleon. The freed nucleons undergo enhanced electron capture while, at the same time, the neutrino opacity of this material has been reduced by about 1/Z because the nuclear coherence has been lost. Thus, as the shock wave approaches the trapping density of $\sim 10^{12}$ g/cm^3, there is a sudden increase in the star's electron neutrino luminosity (sometimes referred to as the deleptonization pulse) to values in excess of 10^{54} ergs/sec. The duration of this pulse, a few milliseconds, is comparable to the time required for shock wave propagation ($v_{\rm shock} \sim 0.2$ c) from the edge of the homologous core ($\sim$ 50 km) to the trapping radius ($\sim$ 100 km). The deleptonization burst carries, perhaps, $3 \cdot 10^{51}$ ergs. Over the next $\sim$ ten seconds the neutron star cools by emitting pairs of neutrinos of all flavors, which carry off the bulk ($\gtrsim$ 99%) of the $\sim 3 \cdot 10^{53}$ ergs released by the collapse. A plot of the neutrino luminosity is given in Fig. 1. (Note the broken time scale.)

Precisely how the mantle of the star is ejected is still a matter of some controversy. Heroic efforts have been invested in numerical codes for modeling the

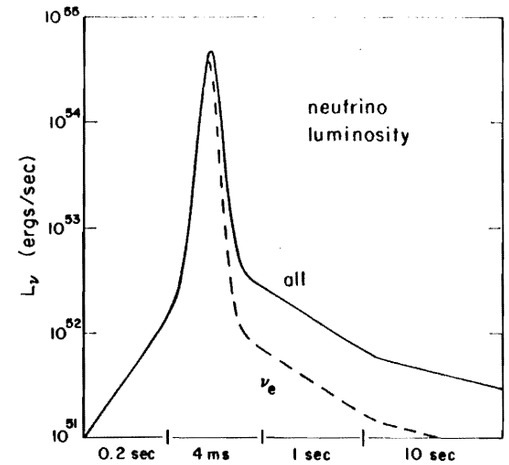

Fig. 1. Schematic representation of the neutrino luminosity of a core-collapse supernova (modified from the results of Ref. [9]). Note the broken scale. The first 0.2 sec corresponds to infall; the 4 ms period represents the time between core bounce and shock breakout at the neutrinosphere; 1 sec is the time from propagation to the edge of the iron core; and 10 sec is the neutron star cooling time.

hydrodynamics of shock wave propagation and the diffusion of neutrinos. While initially the shock wave has an energy far in excess of that needed to eject the mantle, the heating and neutrino losses it experiences in the outer iron core are severe. In the "prompt mechanism" the shock wave survives these losses with enough energy to power a successful explosion.[9] Previously it was thought that this mechanism could be successful in smaller stars ($M \lesssim 15\ M_\odot$), provided the nuclear equation of state is soft, permitting the "rebound" to occur at very high nuclear densities. The inclusion of neutrino "downscattering" mechanisms[11] in the infall stage (inelastic reactions by which superthermal electron neutrinos produced in electron capture can equilibrate with the matter, effectively lowering the average neutrino opacity) reduces the lepton number trapped in the core, which in turn leads to a smaller homologous core, a weaker shock wave, and a larger outer iron core for the shock wave to penetrate. As a result, the prompt model now appears to work only for rather extreme choices of the equation of state. In the "delayed mechanism", the

shock wave stalls after $\sim$ 200 km, going into an acretion mode.[10] It is revived about 0.5 sec later due to neutrino heating of dissociated nucleons behind the shock front. This mechanism is still poorly understood and appears to depend on the detailed modeling of convection that occurs in the hot "bubble" that forms between the neutron star surface and the stalled shock front.[10]

These explosion models have been and are being improved through more realistic treatments of the hydrodynamics, neutrino diffusion, and neutrino interactions. I would like to discuss several theoretical efforts on the last topic that give one some hope that the inclusion of additional physics in the hydrocodes could lead to a more robust mechanism.

I previously mentioned that the infall epoch sets the "initial conditions" for the core bounce, including the trapped lepton number Y_L that governs the size of the homologous core. The star radiates lepton number in the conversion $e^- + p \rightarrow n + \nu_e$. The rate of electron capture varies at T^5, where T is the matter temperature, with T rising as the infall phase proceeds. Furthermore, lepton number losses cease once trapping density is reached, $\rho \sim 10^{12}$ g/cm^3. Thus, as can be seen from Figure 2, a narrow window in density exists where T is sufficiently high but ρ is sufficently low that significant losses in Y_L occur. In fact, the Y_L loss depends critically on processes by which superthermal ν_es produced in electron capture, which have larger coherent cross sections, can scatter into lower energy states, which have an easier time escaping. If one turns off the inelasticities $\nu_e + e \rightarrow \nu_e' + e'$ and $\nu_e + A \rightarrow \nu_e' + A'$, Y_L of the homologous core remains large, ~ 0.41.[12] But either of these inelasticities is sufficiently robust that it can fill the phase space of low-energy unoccupied neutrino states as quickly as they are emptied by diffusion, leading to a final $Y_L \sim 0.38$ (see Fig. 2). Such small values of Y_L imply a smaller homologous core, a thick outer iron core, and great difficulties for the prompt explosion model.

It is therefore important to examine processes that might slow this lepton number loss. One helpful bit of physics emerged from recent calculations of Bruenn and Haxton[12], who rediscovered a result known earlier to Cooperstein and Wambach.[13] Electron capture prior to trapping is conventionally modeled as a heating process: shell model arguments (an $f_{7/2}$ proton becomes an $f_{5/2}$ neutron) suggest that most of the beta decay strength in iron-like isotopes is carried by excited states residing $\sim$3 MeV above the Fermi surface in the daughter nucleus.[14] Thus shell model overlaps are maximized by the capture of an electron on a "cold" parent nucleus, with the production of a "hot" daughter and an escaping neutrino. In fact, the more detailed calculations mentioned above demonstrate that phase space is a much more important factor than the nucleon overlaps in realistic nuclear calculations. Typical decays proceed from highly excited initial nuclear states ($\sim 6kT$) to low lying states ($\sim 1kT$) in the daughter, thereby producing a large energy release.[12] Thus, weak interactions tend to maintain a low-entropy collapse, which of course slows electron capture and contributes to a larger trapped Y_L.

Similarly, Haxton and Lee[15] and Fuller and Meyer[16] have pointed out that neutral current β decay, $A' \rightarrow A + \nu + \bar{\nu}$, with the neutrinos escaping, will cool the star without lepton number loss. This process is not included in present hydrody-

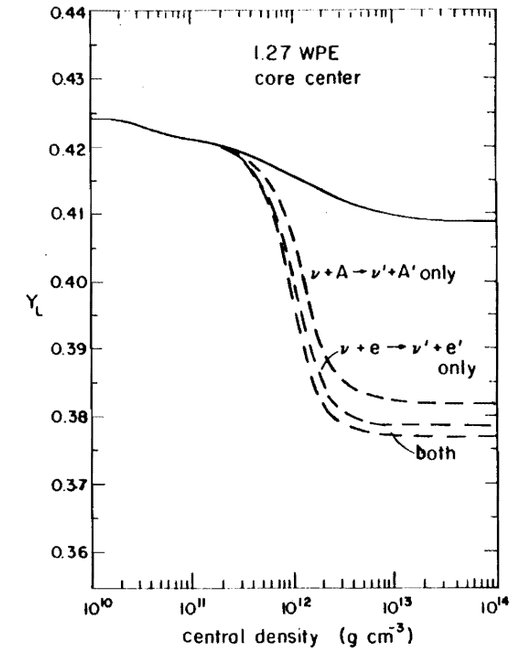

Fig. 2. Trapped lepton number Y_L as a function of density for a moving mass element near the core.[12] The solid curve shows the results for a collapse without neutrino downscattering (i.e., with omission of neutral current scattering off nuclei and of neutrino-electron scattering), while the dashed lines show the effects of turning on the $\nu_e + e \rightarrow \nu_e' + e'$ and $\nu_e + A \rightarrow \nu_e' + A'$ inelasticities.

namic models of the supernova mechanism. In the allowed limit, the rate would be determined by the distribution of M1 strength near the nuclear groundstate. As Y_L drops to $\sim$ 0.41 - 0.42, several odd-mass nuclei (^{51}V, ^{53}Cr, ^{55}Mn) become important species in nuclear equilibrium, comprising about 10% of the matter.[16] These nuclei do have strong low-energy M1 gamma transitions. Recently Horowitz has estimated plasmon enhancements of neutral current beta decay, which are very

important for electric transitions and time-like momentum transfers.[16] Thus it is possible that forbidden transitions (E1, E2) could prove to dominate neutral current β decay rates.

As the proper modeling of electron capture and neutral current β decay will help maintain a lower entropy during collapse, their inclusion in numerical simulations might be important. However, it is not yet known whether the effects substantially alter the evolution.

In the late-stage model, where the shock stalls at a radius of about 200 km, the reheating of the electron and nucleon gas behind the shock front is thought to power the shock wave revival. The heating processes included $\bar{\nu}_e + p \rightarrow n + e^+$ and the quasielastic scattering of neutrinos of any flavor off nucleons, $\nu + N \rightarrow \nu' + N'$, a reaction that transfers only a few MeV, typically. Haxton and Lin have discussed a new process, inelastic neutral current scattering off correlated nucleon pairs, which depends on the second-order polarization insertion at low densities.[17] It will be interesting to see whether this and other neutrino heating mechanisms lead to a more robust delayed mechanism.

4. The Neutrino Process

The analog of this process for nuclei, inelastic scattering of neutrinos, plays an important role in nucleosynthesis, as Woosley, Haxton, and collaborators[18] have recently shown. As this work is published, the description here will be very brief. The importance of these reactions can be seen immediately. Consider the neon shell, which resides at $\sim (1-3) \cdot 10^4$ km from the center of the star. The higher temperature $\nu_\mu s$ and $\nu_\tau s$ have cross sections for inelastic scattering of $\sim 3 \cdot 10^{-41}$ cm^2, with typical energy transfers of $\sim$ 20 MeV. (The Gamow-Teller and various dipole resonances dominate the scattering.) The total neutrino fluence at $2 \cdot 10^4$ km is $\sim 10^{38}$/cm^2. Thus we discover that approximately 0.1% of the neon nuclei are excited above particle breakup by the neutrinos emitted by the cooling neutron star.

The only stable isotope of fluorine, ^{19}F, has a solar abundance relative to Ne of ^{19}F/^{20}Ne $\sim$ 1/3100. The astrophysical site for production of this isotope, however, had been unclear. When excited to $\sim$ 20 MeV, ^{20}Ne* decays into ^{19}F +p ($\sim$ 2/3 branch) and ^{19}Ne + n ($\sim$ 1/3 branch), with ^{19}Ne decaying into ^{19}F. Thus we conclude, from the fluence and cross section given above, that the ^{19}F/^{20}Ne ratio produced in the neon shell, were all the ^{19}F to survive, would be about 8 times the solar ratio. [Thus ^{19}F would be overproduced if Type II supernovae are the principal site for ^{20}Ne production.]

But a key issue is the survival of ^{19}F. In the first 10^{-8} sec after decay into ^{19}Ne + n, the coproduced neutron is absorbed by one of the important neutron poisons, ^{15}O(n,p)^{15}N, ^{19}Ne(n, α)^{16}O, ^{20}Ne(n,γ)^{21}Ne, ^{19}Ne(n,p)^{19}F, etc. The reactions that destroy ^{19}Ne account for $\sim$ 70% of this processing.

In $\sim 10^{-6}$ sec protons produced by ^{20}Ne$^* \rightarrow$ ^{19}F+p are absorbed through ^{15}N(p,α)^{12}C, ^{19}F(p,α)^{16}O, ^{23}Na(p,α)^{20}Ne, etc., with the competition between the latter two depending on the metallicity of the star.

The ^{19}F that survives this processing can still be destroyed by the heating accompanying shock wave passage, ^{19}F(γ,α)^{15}N. The peak explosion temperature spans the range (2.5 - 1.0) $\cdot 10^9$ K over the neon shell ($R \sim (1-3) \cdot 10^4$ km). The survival temperature for ^{19}F is $1.7 \cdot 10^9$ K. Thus only that ^{19}F produced at large R survives the passage of the shock wave.

Table 1. Neutrino process production[19] of ^{19}F relative to ^{20}Ne, normalized to the solar ratio, as a function of the ν_μ/ν_τ temperature. For simplicity, Fermi-Dirac neutrino spectra have been used.

$[^{19}F/^{20}Ne]/[^{19}F/^{20}Ne]_\odot$	T_{ν_μ} (MeV)
0.14	4
0.6	6
1.2	8
1.1	10
1.1	12

The net results for the ^{19}F/^{20}Ne ratio, normalized to the solar abundance, are given in Table 1 as a function of the ν_μ/ν_τ temperature T_{ν_μ}. Excellent agreement with the solar abundance is obtained provided $T_{\nu_\mu} \gtrsim 7$ MeV, a value in good accord with theory. Interesting, although the nuclear cross section rises extremely rapidly with T_{ν_μ}, the production plateaus above 8 MeV. This reflects the competition between the proton poisons ^{19}F and ^{23}Na: if ^{19}F is too abundant, it becomes the dominant capturing nucleus, thereby destroying itself. Thus a rather interesting prediction emerges: the "neutrino process," in the case of ^{19}F, is not really primary, as stars with higher metallicities will be richer in Na and thus more efficient factories for ^{19}F. It follows that the galactic production of ^{19}F should track Na, an interesting nucleosynthetic correlation that may be testable.

Detailed network calculations have been carried out for the neutrino process, with the result that ^{11}B, ^{7}Li, ^{19}F, ^{138}La, and ^{180}Ta all appear to be produced at about their galactic abundance.[19] The production of about 15 other odd-A isotopes is also quite robust, enough so that, within astrophysical uncertainties, the neutrino process could also be the primary source of these nuclei. It thus appears that the neutrino process is an important nucleosynthetic mechanism that contributes, along with explosive burning, to the production of many of the less abundant elements.

5. Conclusions

It is remarkable that the supernova neutrinos Fred Reines has studied on earth also play an important role in the explosion mechanism and in nucleosynthesis. In fact, the abundance of elements like ^{11}B and ^{19}F appear to constrain the flux and temperature of those heavy-flavor neutrinos we did not detect from SN1987A, thereby adding to the information we gleaned from IMB and Kamioka II. It is no accident that simultaneously we are making great strides in understanding the supernova mechanism, unraveling the pattern of the elements, and detecting cosmic fluxes of neutrinos. The unifying theme is the quest to better understand the interactions of neutrinos in matter.

This work was supported in part by the U.S. Department of Energy and by NASA under grant #NAGW 2523.

6. References

1. F. Reines, Ann. Rev. Nucl. Sci. **10** (1960) 1.
2. F. Reines, Proc. Royal Soc. **A301** (1967) 159.
3. K. Hirata, *et al.*, Phys. Rev. Lett. **58** (1987) 1490; R.M. Bionta *et al.*, Phys. Rev. Lett. **58** (1987) 1494.
4. S.P. Mikheyev and A. Yu. Smirnov, Yad. Fiz. **42** (1985) 1441 [Sov. J. Nucl. Phys. **42** (1985) 913] and Nuovo Cimento **9C** (1986) 17; L. Wolfenstein, Phys. Rev. **D17** (1978) 2369 and **D20** (1979) 2634.
5. J.N. Bahcall and R.K. Ulrich, Rev. Mod. Phys. **60** (1988) 297 and private communication.
6. B. Pontecorvo, Zh. Eksp. Teor. Fiz. **36** (1959) 1615 [Sov. Phys. JETP **9** (1959) 1148]; G.M. Gandel'man and V.S. Pinaev, Zh. Eksp. Teor. Fiz. **37** (1959) 1072 [Sov. Phys. JETP **10** (1959) 764]; G. Beaudet, V. Petrosian, and E.E. Salpeter, Ap. J. **150** (1967) 979.
7. G.G. Raffelt and D.S.P. Dearborn, Phys. Rev. **D36** (1987) 2211; M. Fukugita, S. Watamura, and M. Yoshimura, ibid. **26** (1982) 1840; D.A. Dicus *et al.*, ibid. **22** (1980) 839; W.C. Haxton and K.Y. Lee, Phys. Rev. Lett. **20** (1991) 2557.
8. G.G. Raffelt, Phys. Rev. Lett. **64** (1990) 2856.
9. S.A. Colgate and H.J. Johnson, Phys. Rev. Lett. **5** (1960) 573; G.E. Brown, H.A. Bethe, and G. Baym, Nucl. Phys. **A375** (1982) 481; J. Cooperstein, H.A. Bethe, and G.E. Brown, Nucl. Phys. **A429** (1984) 527; E.A. Baron, J. Cooperstein, and S. Kahana, Phys. Rev. Lett. **55** (1985) 126 and Nucl. Phys. **A440** (1985) 744.
10. J.R. Wilson, in *Numerical Astrophysics*, ed. J.M. Centrella, J.M. Leblanc, and R.L. Bowers (Jones and Bartlett, Boston), p. 422; R. Mayle and J.R. Wilson, Ap. J. **334** (1988) 909.
11. S.W. Bruenn, Ap. J. **340** (1989) 955 and **341** (1989) 385; E.S. Myra and S.A. Bludman, Ap. J. **340** (1989) 384.

12. S.W. Bruenn and W.C. Haxton, Ap. J. **376** (1991) 678.
13. J. Cooperstein and J. Wambach, Nucl. Phys. **A420** (1984) 591.
14. H.A. Bethe, G.E. Brown, J. Applegate, and J.M. Lattimer, Nucl. Phys. **A324** (1979) 487; G.M. Fuller, W.A. Fowler, and M.J. Newman, Ap. J. **252** (1982) 715.
15. W.C. Haxton and K.Y. Lee, unpublished; G.M. Fuller and B.S. Meyer, Ap. J. **376** (1991) 701.
16. C. Horowitz, talk presented at CEBAF Symposium for J. Dirk Walecka (April, 1992).
17. W.C. Haxton and W. Lin, unpublished.
18. S.E. Woosley and W.C. Haxton, Nature **334** (1988) 45; S.E. Woosley, D. Hartmann, R.D. Hoffman, and W.C. Haxton, Ap. J. **356** (1990) 272.