Astrophysics and Fundamental Neutrino Interactions

M. Ruderman*

New York University, New York, N. Y.

I. INTRODUCTION

In astrophysical and cosmological speculations, the neutrino is often given an exotic role: a symmetric neutrino-antineutrino fluid has been suggested as constituting a primeval fluid, from whose fluctuation our obviously non symmetric local universe has resulted.[1] The possibility of a present universal neutrino degeneracy of enormous energy density (up to 10^{13} MeV/cc) has been offered as a possible consequence of a hypothetical oscillating universe.[2] More conservatively, the neutrino density in our universe may at least considerably exceed that of the so far observable matter (10^{-29} g/cc) and so control the rate of expansion. Such speculative models are not based upon a detailed knowledge of

---

* On leave from University of California, Berkeley, California

the specific way in which neutrinos interact with other elementary particles. They use mainly the observation that this interaction is so weak that even a large flux of low energy neutrinos (e.g. $10^{11}/cm^2/sec$ with $E \sim 1$ MeV) would not yet have been detected. At present, speculative cosmology has no obvious relation the the theory of weak interactions.

It is only in models of the evolution of stars that details of reactions involving neutrinos play decisive roles in contemporary astrophysics. But here it has largely been a question of applying known or extrapolated β-decay and K-capture rates of various nuclear species to the presumed matter within stellar interiors. For example, the proton cycle within the sun or the carbon cycle which predominates in somewhat hotter main sequence stars result in 3-6 percent of the energy of fusion being radiated as neutrinos. From the sun about $10^{11}\ \nu/cm^2/sec$ are expected at the earths surface. A measurement of their intensity and spectrum could differentiate between various details in models for the solar energy cycle (viz. whether B^8 is formed) but would presumably not tell us anything new about β-decay. Similarly the URCA process of Gamow and Schoenberg may be very important as a means of energy loss from a star whose interior temperature exceeds $10^9\ ^oK$ but it is based upon established β-decay mechanisms.

Our concern here will be the question: what is to be learned from astronomical measurements and astrophysical models about the interaction of neutrino with other elementary particles? In particular, what can be said about neutrino interactions that is not yet known from other sources?

2. ELECTROMAGNETIC PROPERTIES OF NEUTRINO[3]

An immediate but rather trivial consequence of varied astronomical evidence about the age of the sun and similar stars is an upper limit to the charge, magnetic moment, and charge radius which, in some cases, is much better than that which has been determined in terrestrial experiments. The existence of a small charge, moment, or charge radius for the neutrino would imply that neutrino anti-neutrino pairs could be electromagnetically produced; virtual photons could be converted into neutrino-antineutrino pairs. In any process in which electron-positron pairs can be made, neutrino pairs can also be produced but with two significant differences: (a) the charge on the neutrino, if any, is very small next to that of the electron so that the probability of electromagnetically emitting neutrino pairs in any interaction is always very tiny; (b) the neutrino mass, if any, is much smaller than that of its associated lepton so that the threshold for neutrino pair emission is very small and may be zero. The neutrinos, if produced at all, easily escape from the interior of a star without further interaction while other forms of energy transmission (via photons or electrons) are limited by the slow diffusion from the interior to the surface. We can exploit the known long life of our sun (at least 5×10^9 years) to put an upper limit on its energy loss through neutrino pair emission and hence on the neutrino electric charge, moment, and charge radius.

We assume that the neutrino mass is not large compared to one KeV; otherwise there is generally not enough energy to create them in stellar interiors where temperatures are typically 10^7-10^8 $^{\circ}$K. In our estimates we take it to be zero.

Perturbation theory for the quantum electrodynamics of massless charged neutrinos is logarithmically divergent in that approximation in which photons have zero mass and infinite mean free path. But within a plasma, quantized transverse electromagnetic waves have the momentum and energy relation of particles with mass:

$$\omega^2 = \omega_p^{\,2} + k^2c^2 \tag{1}$$

where ω is the frequency and k the wave number of the "photon" and ω_p the plasma frequency:

$$\omega_p^{\,2} = \frac{4\pi ne^2}{m_e} \tag{2}$$

n, e, and m_e are the electron density, charge, and mass respectively. At the core of the sun $n \sim 10^{26}$/cc so that the mass of a solar photon, $\hbar\, \omega_p$, is approximately 400 e.v., greater than twice the experimental upper limit to the ν_e mass. (The mean free path of a photon in stellar matter is typically 1 gm/cm^2, corresponding to a lifetime of more than 10^{-13} sec.; the imaginary part of the "photon" mass is then $\sim 10^{-2}$ e.v. and negligible next to the real part.) A massive "photon" can then spontaneously decay into neutrino pairs. To a very good approximation the quantum electrodynamics of massive photons (transverse plasmons)[4] is exactly the same as conventional quantum electrodynamics, except for Eq. (1), i.e., it is identical to the theory for the transverse components in neutral vector meson theory.

If R_o is the decay rate of a "photon" into neutrino pairs in its rest system, its decay rate when $k \neq 0$ is $R_o \frac{\omega_p}{\omega}$ and the rate at which energy is converted into neutrino pairs is simply $\hbar\omega\, R_o \frac{\omega_p}{\omega} = \hbar\, R_o\, \omega_p$ independent of k. The total rate of neutrino emission per unit mass, ϵ, is then

$$\epsilon = \frac{N\, R_o\, \omega_p \hbar}{\rho} \tag{3}$$

and

$$N = 2 \int \frac{d^3k}{(2\pi)^3} \left[\exp\left(\frac{\hbar\omega}{kT}\right) - 1 \right]^{-1} \tag{4}$$

and ρ the mass density. For the decay rate R_o we have for neutrinos of charge e_ν

$$R_o = \frac{e_\nu^2}{\hbar c} \frac{\omega_p}{6} \tag{5}$$

In the special case of a star with $kT \gg \hbar\, \omega_p$, we have the loss per gram of stellar matter, for the production of pairs,

$$\epsilon = 0.04\left(\frac{e_\nu}{e}\right)^2 \alpha\left(\frac{kT}{\hbar c}\right)^3 \frac{4\pi\ ne^2\hbar}{m\rho} \tag{6}$$

For the solar core we take $T \sim 1.5 \times 10^7$, $n \sim 10^{26}$, $\rho \sim 10^2$. Then

$$\epsilon \sim \left(\frac{e_\nu}{e}\right)^2\ 10^{27}\ \frac{\text{ergs}}{\text{gm-sec}} \tag{7}$$

But the visible light radiated by the sun corresponds to an average energy production of about 2 erg/gm/sec. The energy carried away by neutrinos cannot have been more than a factor of 10 greater than this without greatly shortening the life of the sun on the main sequence. For suppose the sun has been emitting 10 times as much energy in neutrinos as in photons over the past 5×10^9 years. The most efficient source of such energy would be the conversion of H to He. From the solar mass and its H content we can estimate for how long the sun could have produced energy at a rate 10 times the visible rate. This turns out to be less than 10^9 years. We can therefore conclude that the neutrino energy loss cannot be too high, and that

$$\left(\frac{e_\nu}{e}\right)^2 < 10^{-26} \tag{8}$$

or

$$\frac{e_\nu}{e} < 10^{-13} \tag{9}$$

This argument depends crucially on the assumption that ν or $\bar{\nu}$ absorption is negligible. If we did not know from other evidence that the interaction of neutrinos with matter was very weak, the neutrinos might be everywhere in thermal equilibrium and thus carry away an energy from the surface which of necessity would be about the same as that of electromagnetic radiation. With a weak coupling the $\nu,\bar{\nu}$ are emitted directly from the hot core rather than the cooler surface as is the case with stellar light.

The solar magnetic field, energy loss in escaping the sun, and any possible neutrino degeneracy have been shown[3] to have a negligible effect on this result. In a similar way upper limits can also be obtained for the magnetic moment and charge radius. These results are summarized in Table I for the two kinds of neutrino, assuming the mass of the muon neutrino is less than 1 KeV.

TABLE I

Property	ν_e	ν_μ
Charge	[$< 4 \times 10^{-17}e$ from charge conservation] $< 10^{-13}e$ from astrophysics $< 3 \times 10^{-10}e$ from electron-neutrino scattering	$< 10^{-13}e$ from astrophysics, if $M_{\nu_\mu} < 1$ KeV. [$< 3 \times 10^{-5}$ from charge conservation] $< 3 \times 10^{-5}$ from pion production by neutrinos
Magnetic moment (in Bohr magnetons)	$< 10^{-10}$ from astrophysics $< 1.4 \times 10^{-9}$ from neutrino-electron scattering	$< 10^{-10}$ from astrophysics, if $M_{\nu_\mu} < 1$ KeV. $< 10^{-8}$ from pion production by neutrinos
Charge radius (in cm)	$< 4 \times 10^{-15}$ from electron-neutrino scattering $< 4 \times 10^{-14}$ from astrophysics	$< 10^{-15}$ from pion production by neutrinos $< 4 \times 10^{-14}$ from astrophysics if $M_{\nu_\mu} < 1$ KeV.

Summary of the known limits for the electromagnetic interactions of neutrinos.[3]

3. ELECTRON-NEUTRINO INTERACTION

A more interesting question to which stellar observational data may offer a clue is that of the coupling of electron-neutrino pairs as represented, at least phenomenologically, by the interaction

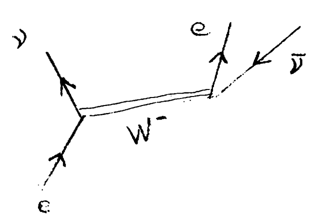
$$H_W = \frac{g}{\sqrt{2}} (\bar{\psi}_e \gamma_\mu (1+\gamma_5) \psi_\nu)(\bar{\psi}_\nu \gamma^\mu (1+\gamma_5) \psi_e) + \text{H.C.} \tag{10a}$$

$$= \frac{-g}{\sqrt{2}} (\bar{\psi}_e \gamma_\mu (1+\gamma_5) \psi_e)(\bar{\psi}_\nu \gamma^\mu (1+\gamma_5) \psi_\nu) + \text{H.C.} \tag{10b}$$

with

$$g = 3.08 \times 10^{-12} \, m_e^{-2} \quad .$$

An effective interaction (10)(for low energy leptons) follows necessarily if the intermediate charged boson exists (and its contribution is not cancelled by a possible appropriately coupled neutral[5,6] counterpart or greatly damped by radiative corrections). Models of universal weak interaction theory which are based upon a product of currents[7] generally contain such a direct $(e\nu)(e\nu)$ coupling.



According to (10b) an electron can radiate a $\nu\bar{\nu}$ pair in part with exactly the same matrix element as that for electromagnetic radiation but with greatly reduced probability. Thus for example, in an atomic transition of energy E, the probability for radiation of a $\nu\bar{\nu}$ pair, $R_{\nu\bar{\nu}}$, relative to that for radiating a photon, R_γ, is

$$\frac{R_{\nu\bar{\nu}}}{R_\gamma} \sim \frac{g^2 E^4}{e^2 \hbar^5 c^5} \sim 10^{-21} \left(\frac{E}{m_e c^2}\right)^4 \tag{11}$$

much too small for direct detection. The probability for emitting a $\nu\bar{\nu}$ pair rises rapidly with energy* (in this case like E^7) but even en electron-positron pair will annihilate into $\nu\bar{\nu}$ rather than

---

* An electron moving with frequency ω in a classical circular orbit of radius R radiates electromagnetic waves, gravitational waves and neutrinos with powers P:

$$P_\gamma \sim \frac{e^2}{c^3} R^2 \omega^4$$

$$P_{GRAV} \sim \frac{GM^2 R^4 \omega^6}{c^5} \quad \text{(quadrupole radiation)}$$

$$P_{\nu\bar{\nu}} \sim \frac{g^2}{\hbar c^8} R^2 \omega^8$$

Neutrino radiation is even weaker than gravitational at very low frequencies. Even for a short period (1 hr) double star (of say H and He respectively so that the center of <u>electron</u> charge moves but not the center of mass) the dipole radiation of neutrinos is negligible even next to the very small quadrupole radiation of gravitational waves.

a photon pair less than once in 10^{22} times. Confirmation of the interaction (10) in a terrestrial experiment would then seem to depend upon the measurement of elastic neutrino (or anti-neutrino)-electron scattering with a predicted cross section $\sigma \sim 8 \times 10^{-45} E$ (E is the laboratory energy in MeV). But astronomical observations when considered together with astrophysical models do seem to offer a possibility of testing this electron-neutrino interaction. For in a normal star the mean free path for neutrinos ($E \lesssim 1$ MeV) is much greater than the stellar radius. All neutrinos that are produced in the stellar interior will escape while photons, with a mean free path equivalent to perhaps a gram of matter, will diffuse slowly and be radiated from the relatively cool stellar surface. Thus the electron-neutrino coupling, if it exists, can be an important and even dominant mechanism for energy loss by a star in certain phases of its evolution. The direct detection of such neutrino pairs from a distant star is not feasible except perhaps from a nearby type II supernova or during a possible subsequent rapid collapse to a neutron star. Indirect evidence comes largely from observations of the distribution of luminosities and surface temperatures of stars and from measured element abundances and the inferences that can be drawn from such data about the stellar interior and the rate of stellar evolution.

4. NEUTRINO EMISSIVITIES

Mechanisms for the radiation of $\nu\bar{\nu}$ pairs by electrons within a star include the following: (see Fig. (1))

(a) $\gamma+e \rightarrow e+\nu+\bar{\nu}$ (photo neutrinos)[9,10,11]

(b) $e^-+e^+ \rightarrow \nu+\bar{\nu}$ (pair annihilation neutrinos)[12,13]

(c) $\Omega \rightarrow \nu+\bar{\nu}$ (plasma neutrinos)[4]

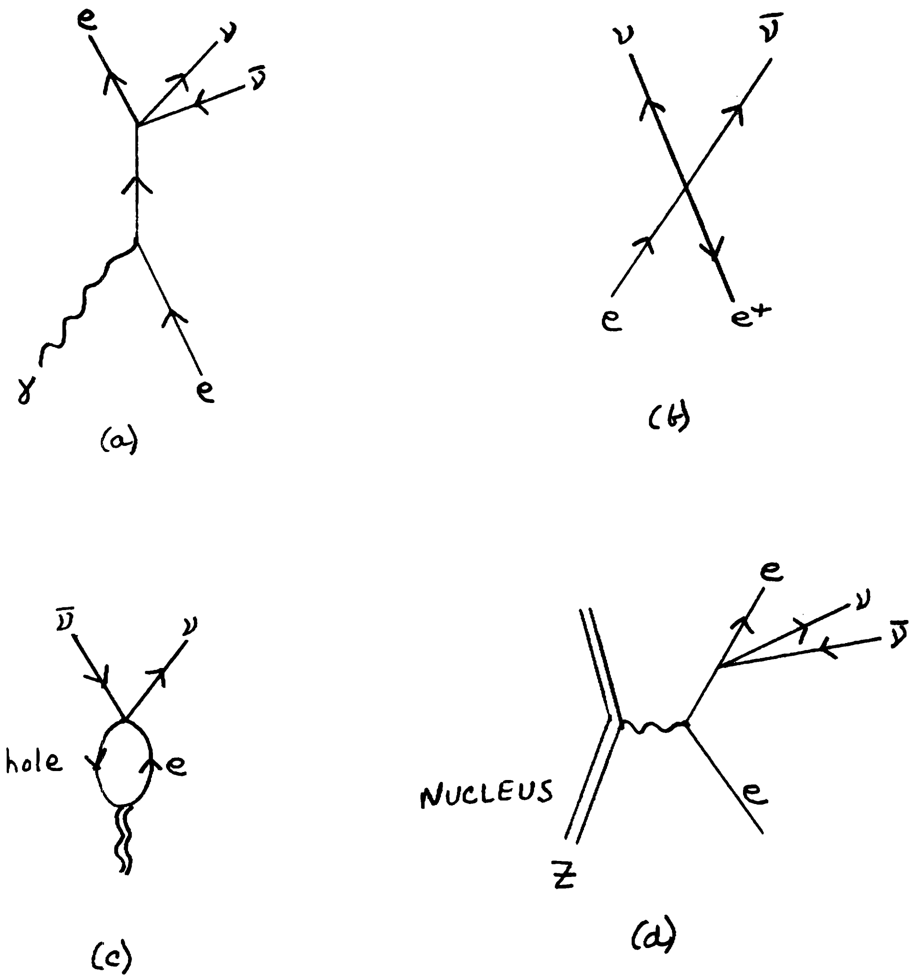
Ω is a collective electromagnetic plus electron wave in the stellar plasma

(d) e+nucleus $\rightarrow$ e+nucleus+$\nu+\bar{\nu}$ (neutrino Bremsstrahlung)[14]

(e) γ+nucleus $\rightarrow$ nucleus+$\nu+\bar{\nu}$ (photonuclear neutrinos)[15,16]

etc.

Process (d) was that offered by Pontecorvo[8] in his original suggestion of the possible importance of such processes in stellar evolutions. Processes (a), (b), and (c) are the dominant ones in those regimes of density and temperature typical of stellar interiors. Figure 2 specifies those domains in which these various neutrino emission processes are prominent. The emissivity q for photonuclear neutrinos[9,11] in fully ionized non-degenerate matter is



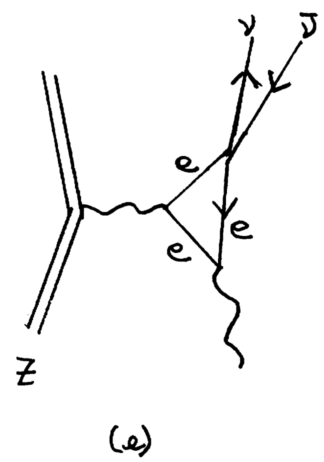


Figure 1: Typical Feynman Diagrams for contributions to neutrino pair emissivity for processes (a)-(e).

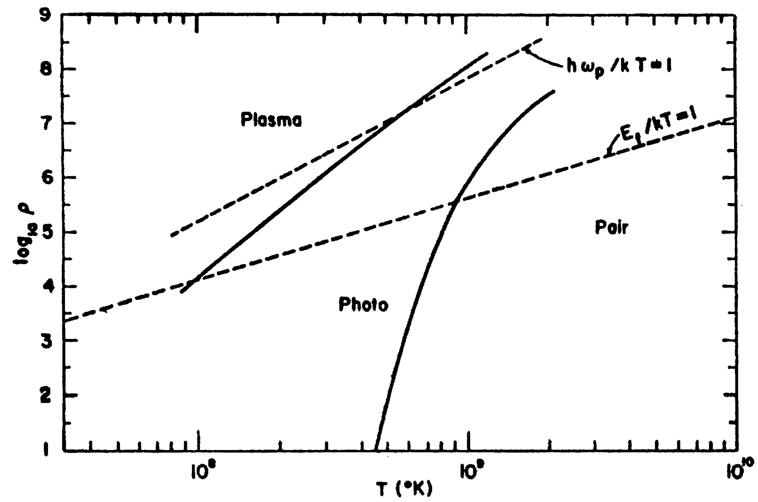


Figure 2: Domains in which different neutrino emission mechanisms dominate.

$E_f \equiv$ Fermi energy

$\omega_p \equiv$ Plasma frequency

Figure from reference 17.

$$q_a \sim 0.6\,(T_8)^8 \text{ ergs/gm/sec} \tag{12}$$

where T_8 is the temperature in multiples of $10^{8}\,^{o}$ K.

In degenerate matter

$$q_a \sim \frac{0.2}{\rho^{1/3}}\,(T_8)^9 \text{ ergs/gm/sec} \tag{13}$$

with ρ the density (gms/cc).

For the pair neutrino emissivity[13] in non-degenerate matter

$$q_b \sim \frac{5\times10^{15}}{\rho} T_8{}^3 \exp(-\frac{119}{T_8}) \text{ ergs/gm/sec}$$

$$T_8 \ll 60 \qquad (14)$$

and

$$q_b \sim \frac{4\times10^{6}}{\rho} (T_8)^9 \text{ ergs/gm/sec}$$

$$T_8 \gg 60 \qquad (15)$$

At $T_8 \sim 50$, $kT = m_e c^2$.

Plasma neutrino emission is significant, if ever, at lower temperatures or higher densities than (a) and (b), and has a more complicated density dependence. Some results and comparisons are given in Figs. 3 and 4.

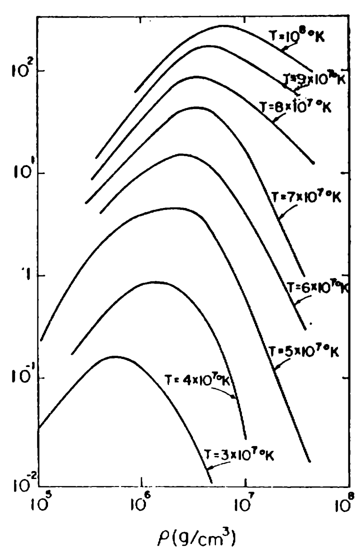


FIG. 3 Emissivity from transverse plasmons q_t as a function of density ρ and temperature.

Figure from reference 4

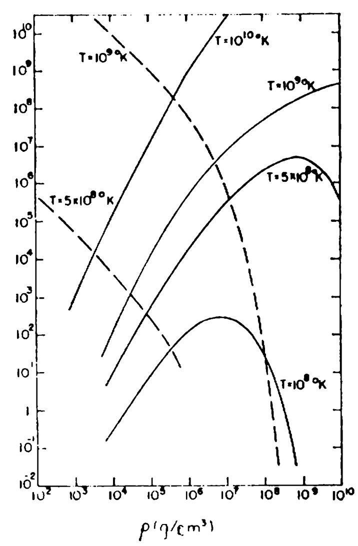


FIG. 4 Comparison of emissivity from transverse plasmons q_t with that from electron-positron pair annihilation (dotted lines), as a function of plasma density and temperature.

Figure from reference 4.

The rates of energy loss from (a), (b), and (c) drop sharply with decreasing temperature. They are quite negligible, for example, in the sun: with a central temperature $T_c \sim 2 \times 10^7 \, ^oK$ a central density $\rho_c \sim 100$ g/cc the neutrino emissivity is infinitesimal relative to the energy production from nuclear burning (~ 2 ergs/gm/sec) and the surface photon emission. But for stars which have evolved off the main sequence with hotter cores, neutrino pair emission can be significent and for $T \gtrsim 6 \times 10^8 \, ^oK$ it would be the controlling energy loss process for a star.

5. STELLAR EVOLUTION

A star is a contracting mass of hot gas in quasi-static equilibrium. The virial theorem then relates the average kinetic energy (temperature) to the average potential energy:

$$k\bar{T}_s \sim \frac{G\, M_s M_z}{R_s} \sim G M_z \, M_s^{3/2} \, \bar{\rho}_s^{\,1/3} \qquad (16)$$

where M_s, $\bar{R}_s$, $\bar{T}_s$, and $\bar{\rho}_s$ are the stellar mass, and the appropriately averaged radius, temperature and density, and M_z is the nuclear mass. (The effects of radiation pressure, magnetic field, and rotation are neglected here.) A hot mass of gas loses energy by radiation. If exothermic nuclear reactions are not present the source of the radiated energy is gravitational contraction; half of the lost gravitational potential energy is radiated away and half is converted to kinetic energy. Thus a star as it radiates is expected to grow hotter and denser roughly according to $\bar{\rho}_s \sim (\bar{T}_s)^3$.

Two effects can interrupt this contraction: (1) Nuclear reactions may serve as the chief source of radiated energy and the contraction may be delayed until an appreciable fraction of the nuclear fuel has been consumed. This is the case of the sun and all main-sequence stars. In general stars spend most of their lives burning fuel and relatively much shorter times in stages of contraction. (2) The core of the star may become degenerate as the Fermi energy increases more rapidly than kT (non relativistically $E_F \sim \rho^{2/3}$). According to Eq. (16) at fixed T, heavier stars have smaller densities and theoretical calculations for the evolution of a heavy star,[18] $M > 10\ M_\odot$, ($M_\odot \equiv$ the solar mass, 2×10^{33} grams) indicate that the density probably never becomes sufficiently high for degeneracy to be important as the star evolves from a hydrogen burning main sequence star to a possible supernova with a central temperature of about 6×10^{9} °K. For lighter stars degeneracy can indeed by significant but it cannot permanently interrupt the stellar contraction if $M > 1.2\ M_\odot$ (the Chandrassekhar limit). For a star with $M < 1.2\ M_\odot$ (or about $4\ M_\odot$ if it is pure hydrogen) the degeneracy pressure can balance the gravitational contraction. It is then possible for radiation to cool the star since it may not be accompanied by an appreciable gravitational contraction. This is thought to be the state of white dwarfs.

Neutrino emission processes can effect the stellar evoluation of light and heavy stars differently. For a heavy star ($M \gtrsim 10\ M_\odot$) it can only speed up the rate of evolution; the additional removal of energy from the stellar interior is accompanied by faster nuclear burning and gravitational contraction. But for a light star with a degenerate stellar core the energy for neutrino emission may come almost solely from the nuclear kinetic energy; the result of neutrino emission may then conceivably be a cooling of the core and a slowing up of the rate of evolution or possibly even the termination of it (if $M < 1.2\ M_\odot$) as the core is prevented from becoming hot enough for subsequent stages of nuclear burning (viz. He $\rightarrow$ C or carbon fusion).

Unfortunately one can only see the stellar surface. To correlate it with conditions within the star and especially within the stellar core requires a reasonably correct model for the evolving star. Then, for appropriately chosen clusters of stars and a model to relate the surface properties to those of the core, the rate of evolution through given interior states can, in principle, be inferred from the number of stars of the appropriate type which are observed. Where evolution is very rapid, few stars should be seen and where it is very slow we should find most stars.

Figure 5 is a summary of the observed distribution of luminosities and surface temperatures for certain clusters of stars within the galaxy. The stars in each individual cluster

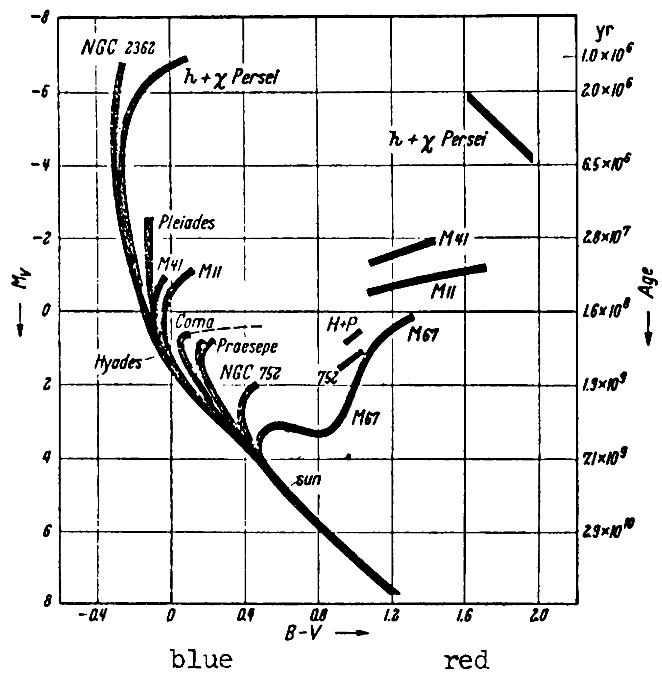


Figure 5: Magnitude vs Color for galactic clusters (surface temperature decreases to the right). The ordinate scale on the right gives the length of time a main sequence star of given M_V is expected to burn before moving off. (from A. Sandage, Astrophys. J. 135, 333 (1962)).

(numbering typically $10^2 \sim 10^3$) are presumed to have been born at about the same time. (An accidental clustering of stars is *a priori* too improbable.) These galactic clusters are composed of so called population I stars, i.e. second generation stars which presumably were formed from gas which contained some of the nuclear products of earlier nucleogenesis. The ordinate $M_V \equiv$ visual magnitude and is semiempirically related to the absolute luminosity. The abscissa B-V is the color scale and measures the surface temperature; increasing B-V corresponds to decreasing temperature. This M_V vs B-V plot is a Hertzsprung-Russell (HR) diagram and it summarizes the empirical evidence for the evolution of the surface of a star: certain parts of the curve for a given cluster are thought to be related to changing surface temperatures and areas during the life of a typical star.

In recent years astrophysicists have made enormous progress in understanding the relation between the surface appearance and the internal state of stars near the main diagonal ("main sequence") in the HR diagram.

Stars spend most of their lives on the main sequence fusing H to He via the p-p cycle or the carbon cycle in slightly hotter cores. Possible neutrino emission via e-ν pair coupling gives an entirely negligible contribution to the energy loss of these main sequence stars.

In heavier stars ($M \gtrsim 10\ M_{\odot}$) H is first burned in a convective core until it is entirely consumed there. The He core begins to collapse and get hotter. H burning then takes place in a shell source surrounding the He core which continues to grow heavier and hotter. Accompanying this core contraction there is an increase in the stellar radius and a reddening of the emitted light. Such stars move off the main sequence to the right. They are evolving toward giants (in size) or super-giants in the case of the already enormous stars near the top of the main sequence. When the He core temperature approaches 10^{8} °K, He begins to burn to C in lighter stars (and to C, O, and Ne in heavier stars with less dense cores). When the He at the center is exhausted the C core begins to collapse and heat up until the temperature is sufficient (~ 7×10^{8} °K) for C burning. At this stage one expects[18] a burning convective C core surrounded by a He-burning shell and an enormous convective envelope composed mainly of H. (In the model calculations mixing between regions of different composition is assumed negligible even when a highly convecting region is adjacent to one of different composition!) Finally the core grows still hotter until all the nuclear species of the core are burned to Fe. Then according to Hoyle and Fowler,[19] at a central temperature of around 7×10^{9} °K the core of a heavy star collapses rather suddenly (seconds) as Fe is converted to He in an endothermic reaction.

The star is no longer in quasistatic equilibrium and explodes in a spectacular type II supernova which puts some fraction of the nuclei which have been formed in its interior back into space to be incorporated into future generations of as yet unborn stars (population I).

Lighter and thus denser stars differ from the heavier ones in that the core may become degenerate. Then the ignition of nuclear fuel can raise the temperature without substantially increasing the pressure and the fuel will burn explosively until the temperature rises sufficiently to relieve the degeneracy, increase the pressure and expand the core. This is known as the "flash".

Very light stars ($M \lesssim 1.2\ M_{\odot}$) can be kept from contracting by the pressure of the degenerate core electrons and may begin to cool to white dwarfs before the core has evolved very far. In Fig. 6 some typical evolutionary paths are shown for the central temperature T_C and central density ρ_C for a heavy, a medium, and a light star as given by Hayashi, Hōshi, and Sugimoto.[18] In Fig. 7 the corresponding calculated luminosities and surface temperatures T_e are shown. For the star of $15.6\ M_{\odot}$ the computed time scales are given in Table II.

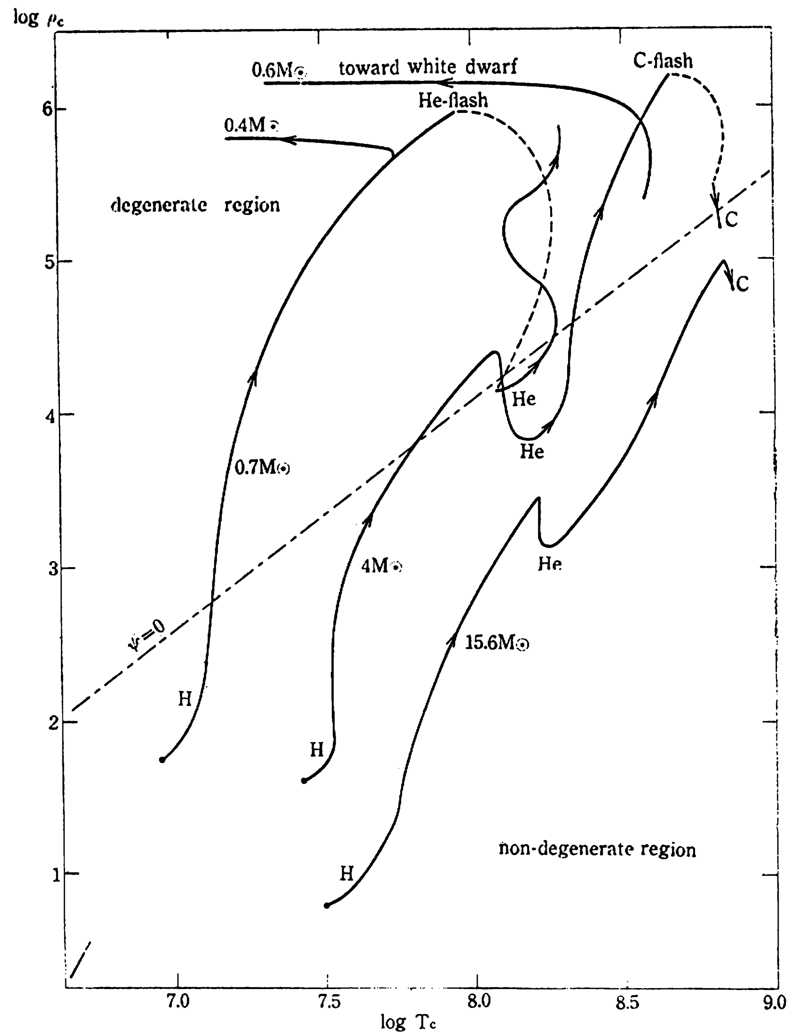


Figure 6: The central temperatures and densities of evolving stars according to Hayashi, Hōshi, and Sugimoto. The line ψ = 0 separates degenerate from non-degenerate cores.

(Figure from reference 18.)

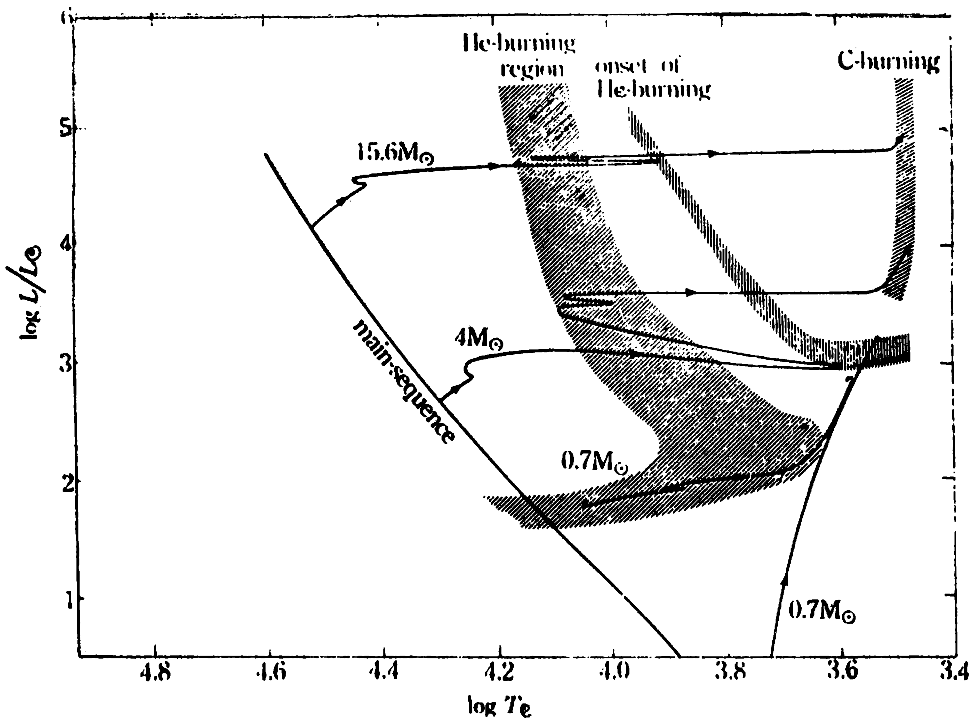


Figure 7: Evolution of the stellar surface corresponding to the interior conditions of Fig. 6 according to Hayashi, Hōshi and Sugimoto. Stars spend most of their lives in the hatched regions.

(Figure from reference 18.)

TABLE II

Computed lifetimes in various state of stellar evolution for a star of 15.6 $M_{\odot}$ as given by Hayashi, Hōshi, and Sugimoto.[18] The C-burning core is taken to be initially one half C and one half O, Ne, Etc. The later phases are computed on the *ad hoc* assumption that $\log(L_{core}/L_{\odot}) = 4.7$ throughout the red supergiant phase.

Lifetime in 10^5 years		
	No photo $\nu\bar{\nu}$	With photo $\nu\bar{\nu}$
Core H-Burning	157	157
Contraction	0.7	0.7
Core He Burning	11.6	11.6
Contraction	0.5	0.25
Core C-Burning	2.3	~ 0.1
Later Phases	$\lesssim$ 6	0.0

The inclusion of the photo neutrino process greatly accelerates the evolution. If it does not exist the model of Hayashi et al. predicts comparable lifetimes in the He burning phase (blue supergiant) and in the C-burning and later phases (red supergiant). If it does exist the evolution through the red supergiant phase for a star of 15.6 $M_\odot$ is predicted by them to be so rapid that one almost never should observe such a star. To compare these predictions with observations it is significant that almost the entire lifetime of such a star is spent on the main sequence. After it leaves it moves to the right very rapidly. In addition the lifetime τ on the main sequence is a very sensitive function of the stellar luminosity and mass

$$\tau \sim 10^{10} \frac{L_\odot}{L} \frac{M}{M_\odot} \tag{17}$$

$$\frac{L}{L_\odot} \sim \left(\frac{M}{M_\odot}\right)^4 \tag{18}$$

Thus heavier stars burn more brightly and leave the main sequence much more quickly than somewhat lighter ones. Then the interpretation of the evolution of the stars in the double cluster h + χ Persei of Fig. 5 is as follows: these stars were born (i.e. started burning on the main sequence) around about 7×10^6 years ago. Stars with an initial $M_V \gtrsim -3$, corresponding to a mass much above about 16 $M_\odot$, have already evolved off the main sequence and through all the stages of Table II. Stars of mass less than about 15 $M_\odot$, $M_V \lesssim -3$ have not yet evolved off the main sequence. Then the stars of the blue supergiant branch ($-0.2 < B-V < 0.0$) and of the red supergiant branch ($1.6 < B-V < 2.0$) correspond to stars of about 15.6 $M_\odot$; the number of stars in a given region of the H-R diagram is expected to be proportional to the lifetime of a star of 15.6 $M_\odot$ in the region.

Figure 8 shows a computed evolutionary path of the 15.6 $M_\odot$ star computed by Hayashi et al. (shown in Fig. 7) on the HR diagram. The dots are a sample of the stars in h and χ Persei. The ratio of the number of presumed He burning stars (blue supergiants right after C) to presumed stars in C-burning and later stages (red supergiants between e and f) is about 1.5. This is in gross disagreement with the model which led to Table II if photoneutrino emission exists with the strength suggested by the Universal Fermi Interaction. In this latter case the expected ratio should be greater than 20.

If the model of Hayashi et al. for the star of 16.5 $M_\odot$ is correct and if the observed red supergiants are highly evolved stages of such a star then indeed the $(e\nu)(e\nu)$ interaction could not be as large as conjectured as these authors point out. But a variety of circumstances make the argument

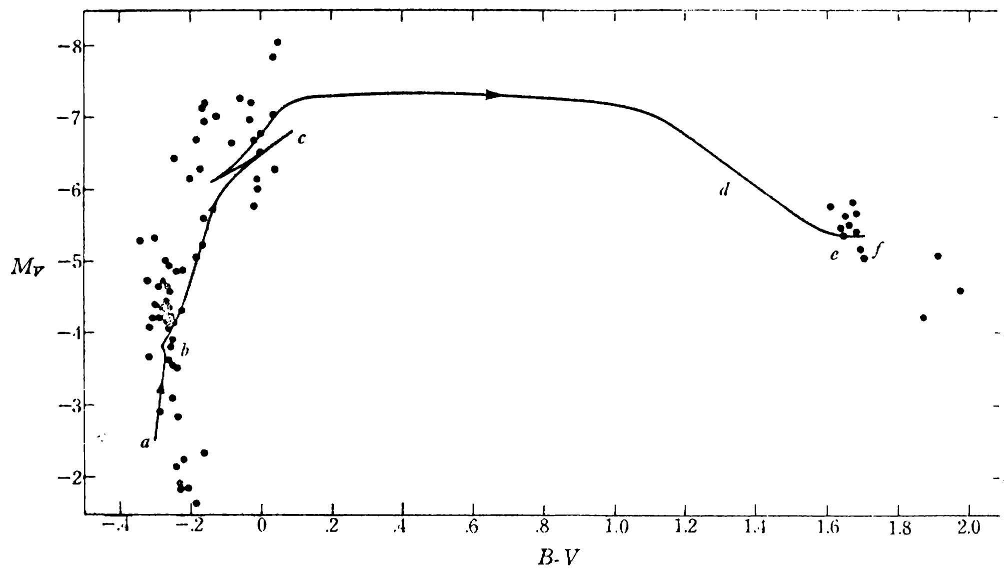


Fig. 8 Evolutionary track of a star of $15.6 M_\odot$ superposed on the color-magnitude diagrams of *h* and χ Persei. Segments of the track correspond to the phases: a-b, hydrogen depletion in the core; b-c, contracting helium core; c-d, helium depletion in the core; d-e, contracting carbon-oxygen core; and e-f carbon burning in the core.

Figure from Hayashi <u>et</u> <u>al</u>., references 18 and 20.

less than compelling at this time:

(a) The results for the time for C-burning depend rather critically upon the core temperature T_c. According to Reeves[17] C at a central carbon density $\rho_c \sim 10^5$ g/cc will burn with $T_c \sim 6.5 \times 10^{8}$ °K (to match the photoneutrino emission) for about 8×10^5 years. This does not differ too greatly from the model results[18] of ρ_c (carbon) $\sim 5 \times 10^4$ g/cc and $T_c \sim 7.7 \times 10^{8}$ °K when C-burning begins. Thus a small change in parameters can greatly lengthen the C-burning phase.* According to Hoyle and Fowler[21] uncertainties involved in computations with highly convective H, He, or C-burning cores and in the deep outer convection zone are sufficient to allow results different from those of Hayashi et al. with somewhat altered but plausible assumptions.

(b) In the relatively undense core of a heavy star of 16 $M_\odot$, C^{12} may be almost completely converted to O^{16} by $C^{12}(\alpha,\gamma)O^{16}$ as it is formed from He burning. According to Fowler this is expected. O^{16} ignites at a higher temperature than C^{12} and if the red supergiants of h and X Persei were burning O^{16} their lifetime would be negligible because of pair neutrino radiation. But even without such neutrino processes, O^{16} burning is probably sufficiently rapid that the red supergiant phase would be very short; an explanation of the number of red supergiants would not seem to follow from the model of Hayashi et al.[21] in either case. Even observationally the relative numbers of red and blue supergiants that are really to be associated with the double Persei cluster are not definitively determined.

---

* Almost all of the observed luminosity can be due to the shell of burning He around the core and may be significantly greater thanthe energy from C burning or gravitational contraction. Thus the luminosity is not necessarily a measure of the rate of C-burning.

(c) The effect of neutrino emission has been estimated as if it were a small perturbation on the state of the stellar core calculated without such emission. Chiu has argued[22] that the inclusion of neutrino effects ab initio may greatly effect the resultant model and alter the present conclusions.

A variety of other effects of possible neutrino pair emission on the lifetime of stellar states have been estimated but so far these cannot be compared to data in a crucial way because of a paucity of observational data or a lack of confidence in quantitative consequences of a particular stellar model calculation.

For example, a star of mass $4\ M_\odot$ will have its C-flash delayed for perhaps an extra 10^6 years[18] but this is about the same as the time for contraction of the He core to begin with. A first estimate of the effect of plasma neutrino emission in possibly completely preventing the He flash in a star of around one $M_\odot$ indicated that it probably would not but the calculation was rather rough.[23]

A significant effect of neutrino pair emission exists for the cooling of a hot white Drawf of mass $0.6\ M_\odot$ with $\log \rho_C \sim 6$ is expected to take over 10^8 years to cool from $T_C \sim 10^{8\,o}K$ to $T_C \sim 5 \times 10^{7\,o}K$ if only surface radiation is possible; its visible luminosity when $T_C \sim 5 \times 10^{7\,o}K$ is only about 1/30 that of the sun. But its neutrino emission is much greater than $L_\odot$ and its lifetime for cooling is thereby reduced from 10^8 to only 10^6 years. Thus one should not expect to see many such hot white drawfs but their visible luminosity is so small that they would probably not be observed even if present in considerable numbers.[24] When the central temperature drops to about 2×10^{7} °K, the neutrino luminosity of a white drawf is negligible next to its visible luminosity.

6. PRE-SUPERNOVA STARS AND NUCLEAR ABUNDANCES

At core temperatures above about 8×10^{8} °K neutrino emission from the annihilation of electrons and positrons (or photo and plasma neutrinos if the density is very high) will certainly control the energy loss rate of a star if the $(e\nu)(e\nu)$ interaction of Eq. 10 exists. One result is to cause a heavy star to evolve extremely rapidly as the core contracts and grows hotter until $T_C \sim 7 \times 10^{9}$ °K when the Fe → He transition is expected to result in a catastrophic core implosion and the star presumably becomes a type II supernova. An estimate of the rate of evolution[25] is given in Table III for $\rho \sim 10^6$ g/cc. (For $\rho \sim 10^8$ g/cc the times are more than 10^2 longer.)

Those $\nu\bar{\nu}$ emission rates may be compared to those expected from the URCA process[26] in which a steady state is achieved between

TABLE III

Neutrino pair emission rates from matter of density $\rho = 10^6$ g/cc and the time scale for core evolution. The virial theorem in the form of Eq. 16 is assumed so that the gain in kinetic energy equals the energy lost. Nuclear burning of Mg, Si, Ne, etc. to Fe does not particularly alter the time scale

Temperature (10^9 °K)	1	2	3	4	5
$\nu\bar{\nu}$ emission (ergs/gm/sec)	10^7	10^{12}	10^{14}	10^{15}	10^{16}
time for increase in temperature of 10^9 °K (sec)	10^{10}	10^5	10^3	10^2	10

β-decay and capture of energetic electrons associated with these high temperatures. Cameron estimates,[27] for $\rho = 10^6$ gm/cc, less than 10^6 ergs/gm/sec at 2×10^9 °K, 10^7 ergs/gm/sec at 2.5×10^9 °K increasing to about 10^{11} ergs/gm/sec at 5×10^9 °K. Then the time scale for the evolution of the core through stages with $T \gtrsim 4 \times 10^9$ °K would be expected to be many years, without the electron neutrino pair coupling compared to $10^2 \sim 10^4$ seconds (for $\rho = 10^6 \sim 10^8$ gm/cc) if it is present. (The neutrino pair emission rate at high temperatures becomes so great that the core will begin to shrink faster than the envelope can follow and the neutrino emission all by itself will result in a catastrophic core implosion if the Fe $\rightarrow$ He transition does not occur first at somewhat lower temperatures.[28])

If the heavy elements in the universe are the result of nucleosynthesis within stars[29] which have subsequently exploded (supernova) then their relative abundances may offer a clue to the conditions and, more significantly for us, the time scale during which they were formed. In particular, the abundances of the elements around Fe^{56} are very suggestive of a thermal equilibrium distribution[30,29] among these elements at a very high temperature $T \sim 3 \times 10^9$ °K and density $\rho \sim 10^5$ or 10^6 g/cc. The equilibrium however can be of two types. Equilibrium among nuclear species occurs very rapidly under these conditions (a few seconds) in so far as it consists of the distribution of a fixed number of neutrons and protons among various nuclei. But if the high temperature equilibrium is maintained for a long time the β-decay rate equals that for K-capture plus positron decay. The resulting mixture of nucleons has a larger fraction of neutrons and the nuclear abundances among isotopes will differ in detail from that of the shorter time scale. Figure 9 is an estimate of the expected abundance distribution for "long" time equilibrium and Fig. 10 shows an expected distribution for "short" time equilibrium compared to observed solar abundances (The sun, a population I star, is assumed to have a distribution of heav elements created in the interior of earlier generations, i.e. population II, stars.) The magnitudes of the calculated abundance ratios depends sensitively upon temperature but which of two iostopes is the major one generally does not. The small observed ratios Fe^{58}/Fe^{57}, Cr^{54}/Cr^{53}, Ni^{62}/Ni^{61} do not obtain for the "long" time scale equilibrium where protons can be converted to neutrons via weak interaction processes. Rather the data argue for the "short" time scale equilibrium for the cooking of the Fe-peak elements. The scale of "short" and "long" depends upon which isotopes contribute to the n-p conversion and their abundances. According to Fowler[31] the equilibrium without appreciable $p \rightarrow n$ conversion necessitates a time scale of no more than 10^3-10^4 sec. Then if the elements around F_e were formed near the outer part of the core of a pre-supernova (II) star and

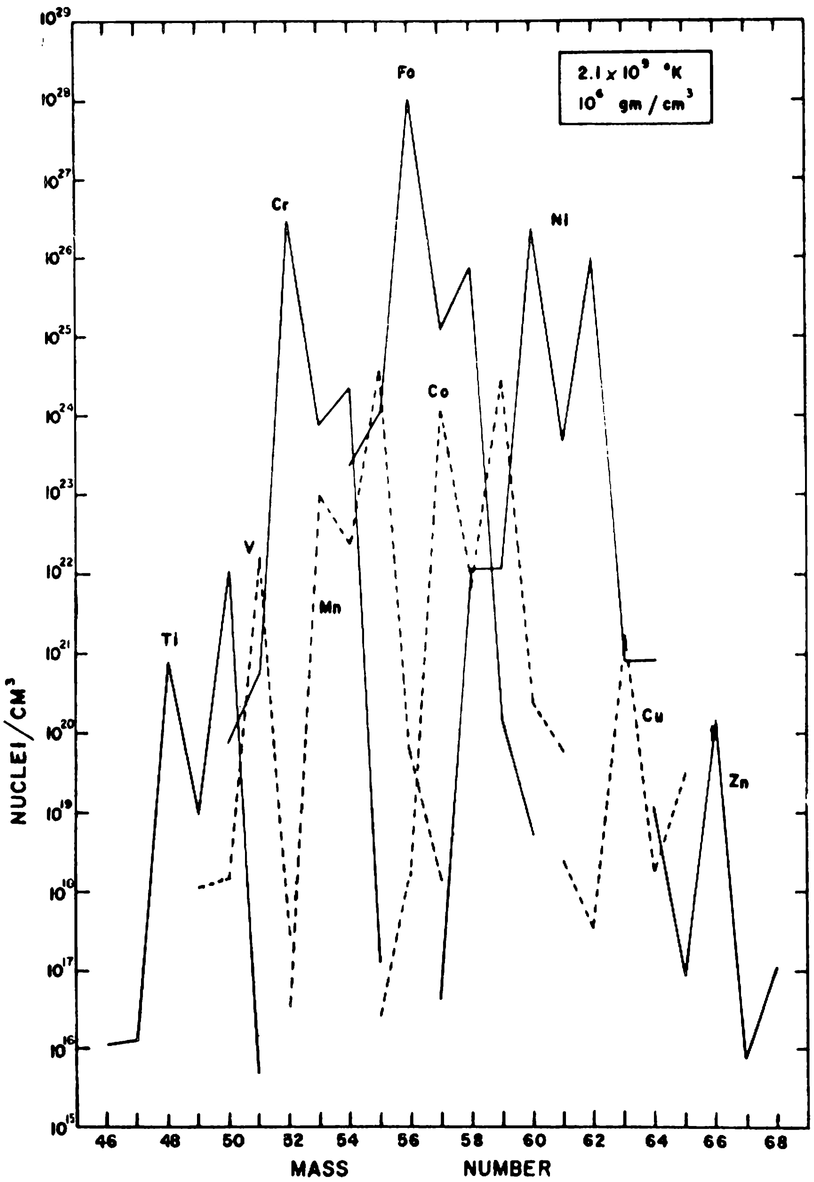
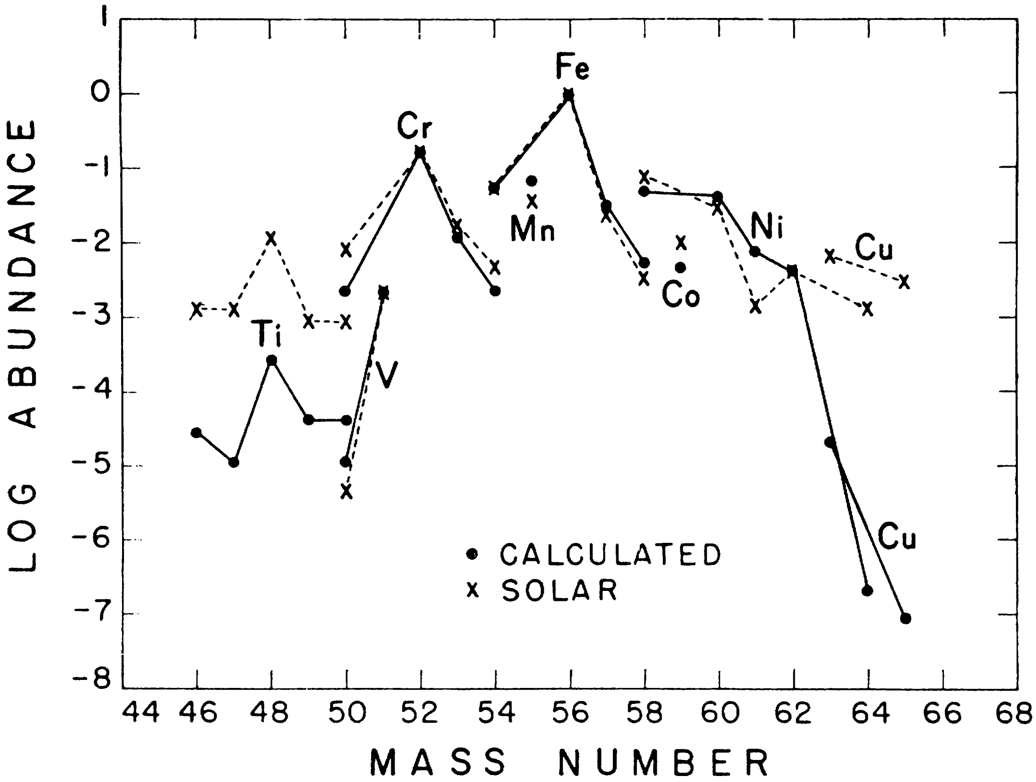


Figure 9: Distribution of expected nuclear abundances for "long" time scale equilibrium as given by A.G.W. Cameron.[27] The ratio of free neutrons to protons turns out to be 6×10^{-3}.



Nuclear statistical equilibrium in the iron peak region as calculated by Burbidge, Burbidge, Fowler, and Hoyle and compared to solar abundances.

Figure 10: The ratio of free neutrons to protons is fixed *a priori* at 3×10^{-3} with $T = 3.78 \times 10^{9}$°K corresponding to $\rho \sim 10^{5}$ g/cc.

(Figure from B^2FH, reference 29).

subsequently ejected in the supernova explosion with a frozen equilibrium distribution, something must enormously accelerate the core evolution from $T \sim 4 \times 10^9$ °K to implosion. The $(e\nu)(e\nu)$ interaction fills this role comfortably. Without it the time for this evolution should be measured in years.

However, other models for the explanation of the iron peak abundances assume that the equilibrium is finally determined only after the core implosion, perhaps after a neutron core has been formed,[27] or even that the elements were formed before the stars.

Thus neither the argument based upon stellar counts and models or that based upon Fe-peak abundances give an unequivocal answer to the question of the weak $(e\nu)(e\nu)$ interaction. The answers they do give are opposite. Perhaps the estimate based upon abundances depends less on detailed model calculations and so may carry more weight but the answer is certainly not conclusive.

In the literature of neutrinos, their weak interactions, and supernova, there is a peculiar abundance of quotations.[2,25,29,32] It is not unfitting then, to recall

"...and God hath chosen the weak things of the world to confound the things that are mighty."

I Corinthians

NOTE

There exists at present a remote possibility for a direct confirmation of the possible enormous $\nu\bar{\nu}$ flux coming from a supernova.

With the Cowan-Reines 10^3 liter scintillator, a fission $\bar{\nu}$ flux of 10^{13} cm^2/sec gave a signal of 3×10^{-2} events/sec, about one fifth background. The pre supernova $\nu\bar{\nu}$ emission from a supernova in our galaxy at say 10^3 light years, would not be detectable with such equipment. But after the core implosion has begun, an enormous $\nu,\bar{\nu}$ burst would be expected within a fraction of a second. For a heavy star the core mass is much greater than the Chandrasekhar limit of 1.2 $M_{\odot}$; thus the catastrophic core implosion once begun will stop only if the temperature grows sufficiently so that thermal pressure stops the collapse (or nuclear densities are reached). But, as has been emphasized by Chiu,[22] the enormous $\nu\bar{\nu}$ emission can remove thermal kinetic energy faster than it can be fed in from the gravitational energy of the collapsing core. The thermal pressure within the core will not appreciably oppose the collapse until the dense star begins to become opaque to ν_e and $\bar{\nu}_e$. Typically more than one percent of the entire rest mass energy of the core could be radiated as $\nu_e,\bar{\nu}_e$ (and more as $\nu_\mu,\bar{\nu}_\mu$) with an energy of at least 10 MeV (the Fermi energy for electrons when $\rho \sim 10^{11}$ g/cc). Then a collapsing supernova core of a few stellar masses within our galaxy at say 10^3 light years, may give hundreds of counts within one second in the Cowan-Reines apparatus if the $(e\nu)$ pair coupling exists. Unfortunately, supernova explosions in the galaxy occur at intervals of hundreds of years.

REFERENCES

1. See, for example, B. Pontecorvo, Soviet Physics, USPEKHI 6, 1 (1963) which contains further references.
2. S. Weinberg, Nuovo cimento (1963).
3. In this section we follow closely J. Bernstein, G. Feinberg, and M. Ruderman, Phys. Rev., to be published.
4. B. Adams, M. Ruderman, and C. Woo, Phys. Rev. 129, 1383 (1963).
5. B. Pontecorvo, Phys. Rev. Letters 1, 287 (1962).
6. S. Bludman, Nuovo cimento 2, 433 (1958).
7. R.P. Feynman and M. Gell-Mann, Phys. Rev. 109, 193 (1958)
8. B. Pontecorvo, Soviet Physics, JETP 9, 1148 (1960).
9. V. Ritus, Soviet Physics, JETP 14, 915 (1962).
10. M. Ida and M. Vahara, (unpublished).
11. H. Chiu and R. Stabler, Phys. Rev. 122, 1317 (1961).
12. H. Chiu and P. Morrison, Phys. Rev. Letters 5, 573 (1960)
13. H. Chiu, Phys. Rev. 123, 1040 (1961).
14. G. Gandel'man and V. Pineau, Soviet Physics, JETP 10, 764 (1960).
15. L. Rosenberg, Phys. Rev. (in press).
16. S. Matinyan and N. Tsilosani, Soviet Physics JETP 14, 119 (1962).
17. H. Reeves, Astrophys. J., 138, 79 (1963).
18. See, for example, C. Hayashi, R. Hōshi, and D. Sugimoto, Supplement to Progr. Theoret. Phys. (Kyoto) No. 22 (1962)
19. F. Hoyle and W.A. Fowler, Astrophys. J. 132, 565 (1960). However, the integrated neutrino flux from all the supernovas within 10^{10} light years (perhaps one per second) may give a measurable cosmic background.

20. C. Hayashi and R.C. Cameron, Astrophys. J. 136, 166 (1962).

21. W.A. Fowler, private communication and to be published.

22. H. Chiu, private communication.

23. H. Chiu, Astrophys. J. 137, 343 (1963).

24. For the effect of neutrino pair emission on the rate of contraction or the cooling of an inert carbon core when $T_C > 10^8$ °K see R. Strothers, Astrophys. J. 137, 770 (1963).

25. Compare also H. Chiu, Ann. Phys. 16, 321 (1961).

26. G. Gamow and M. Schoenberg, Phys. Rev. 59, 539 (1941).

27. A.G.W. Cameron, Mémoires Soc. R. Sc. Liége III (1960). cf. also H. Chiu, Ann. Phys. 15, 1 (1961).

28. H. Chiu and H. Fuller, preprint (1962).

29. E. Burbidge, G. Burbidge, W.A. Fowler, F. Hoyle, Revs. Modern Phys. 29, 547 (1957).

30. F. Hoyle, Astrophys. J. Supplement 1, 121 (1954).

31. Invited paper at Washington Meeting of The American Physical Society, April (1963).

32. H. Chiu, preprint (1963).