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CAN A CLOSURE MASS NEUTRINO HELP SOLVE THE SUPERNOVA SHOCK REHEATING PROBLEM?

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ABSTRACT

We point out that a ν_μ or ν_τ neutrino with a cosmologically significant mass (10–100 eV) and a small mixing (vacuum mixing angle $\theta > 10^{-4}$) with a light ν_e would result in a matter-enhanced Mikheyev-Smirnov-Wolfenstein resonant transformation between these species in a region above the neutrino sphere but below the stalled shock during the reheating phase of a Type II supernova. The neutrino heating behind the shock is due to charged-current ν_e and $\bar{\nu}_e$ captures. Since the ν_μ and ν_τ have considerably higher average energies than do the ν_e, neutrino flavor mixing would result in higher effective ν_e energies behind the shock and a concomitant increase in the heating rate. Our numerical calculations suggest that this effect results in a 60% increase in the supernova explosion energy, possibly helping to solve the energy problem of delayed-mechanism supernova models.

Subject headings: elementary particles — shock waves — supernovae: general

1. INTRODUCTION

In this paper we examine the role that flavor-changing matter-enhanced neutrino oscillations might have in the process of reviving an otherwise dying postbounce supernova shock. Ever since the work of Colgate & White (1966) researchers have sought to explain the mechanism of Type II supernova explosions in terms of the transfer of energy or momentum from the core to the mantle by neutrinos. The motivation for this stems in part from the dominant role played by neutrinos in the energetics of core collapse; where, for example, 95% or more of the gravitational binding energy ($\approx 10^{53}$ ergs) released when the presupernova Fe core drops to a neutron star is transformed into ν_e, ν_μ, ν_τ and associated antineutrino seas. Neutrinos are trapped in the low-entropy infall phase of collapse (duration < 1 s) but diffuse out of the core after bounce on a time scale ~ 10 s (Mazurek 1975; Sato 1975; cf. Arnett 1977 and Bethe et al. 1979 for an overview of the supernova core collapse problem).

Hydrodynamic bounce of the core generates a shock which begins to move out but which suffers energy loss as nuclei are photodisintegrated in the higher entropy postshock environment. In some numerical calculations the shock remains viable and explodes the star with an energy of $\approx 10^{51}$ ergs, while in other studies (especially for massive stars $M \geq 15\ M_\odot$) the shock stalls and evolves into an accretion shock (Cooperstein & Baron 1990; see also Wilson 1982; Baron, Cooperstein, & Kahana 1985; Burrows & Lattimer 1987). In the latter case it has been hoped that neutrinos emitted from a "neutrino sphere" near the edge of the proto–neutron star can reenergize the shock (Wilson 1982; Lattimer & Burrows 1984; Bethe & Wilson 1985, hereafter BW85). The energy of the resulting supernova explosion in these "late-time," or delayed-mechanism, models is at best low, however, being less than $\sim 10^{51}$ ergs (BW85; Mayle 1990; but note also that some recent calculations give energies near 10^{51} ergs).

We show here how neutrino flavor transformations between ν_e and either ν_μ or ν_τ which occur above the neutrino sphere but behind the shock could aid the shock reheating process and result in a delayed-mechanism explosion with an energy of $\geq 10^{51}$ ergs. Interestingly, a neutrino mass level crossing in this region would correspond to vacuum neutrino mass-squared differences of $\approx 10^2$–10^4 eV^2, which could be realized with a closure mass ν_μ or ν_τ. We discuss neutrino oscillations in the supernova environment in § 2, delayed mechanism shock reheating in § 3, modifications in the reheating process effected by neutrino oscillations in § 4, observational constraints on neutrino oscillations in § 5, and conclusions in § 6.

2. NEUTRINO OSCILLATIONS IN DENSE MATTER AND SUPERNOVAE

In this section we will describe the conditions required to get a neutrino mass level crossing in the critical reheating region above the neutrino sphere but below the shock in a nascent Type II supernova. We will argue that $\nu\nu$-neutral current-exchange scattering can usually be neglected relative to νe-exchange scattering in determining neutrino effective masses in the reheating region. The adiabaticity of resonant neutrino flavor transformations in this region will be discussed. We will describe how we follow the evolution of the neutrino distribution functions when neutrino flavor transformations are important.

Matter-enhanced neutrino oscillations, especially as regards the solar neutrino problem, have been extensively considered (cf. Wolfenstein 1978, 1979; Mikheyev & Smirnov 1985; Bethe 1986; Haxton 1986). Resonant neutrino oscillations in supernova cores were treated by Fuller et al. (1987, hereafter FMWS87; see also Fukugita et al. 1988), where it was pointed out that $\nu\nu$-neutral current-exchange scattering can become an important source of effective mass difference between ν_e and either ν_μ or ν_τ. FMWS87 suggested that neutrino oscillations above the neutrino sphere but below the stalled shock could be important for the shock energetics. We expect a matter-enhanced level crossing between ν_e and ν_x (where, hereafter, ν_x

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represents either ν_μ or ν_τ) if the vacuum masses obey the condition $m_{\nu_x} > m_{\nu_e}$. In this case there are no transformations among the antineutrinos. FMWS87 showed that as a ν_e propagates through the stellar material, its forward scattering on the ambient lepton seas via e^- or ν_e exchange generates an effective potential (FMWS87, eq. [2b])

$$A^\mu = (\phi, A) = \frac{G_F}{\sqrt{2}} [\bar{\Psi}\gamma^\mu(1 - \gamma_5)\Psi] , \quad (1a)$$

where Ψ is the target electron or neutrino spinor and G_F is the Fermi constant. This effective potential is equivalent to an effective mass contribution for the propagating ν_e of

$$m_{\rm eff}^2 = 2E_{\nu_e}\langle\phi\rangle - 2\langle \boldsymbol{p}_{\nu_e} \cdot \boldsymbol{A}\rangle + \langle A^2\rangle - \langle\phi^2\rangle , \quad (1b)$$

where the brackets represent an average over the ambient electron or neutrino distribution functions and where $\boldsymbol{p}_{\nu_e}$ and E_{ν_e} are the 3-momentum and energy, respectively, of the propagating neutrino. Values and relative magnitudes for A and ϕ in the context of supernova collapse conditions are discussed at length in FMWS87.

In FMWS87 isotropic lepton distributions were considered so that the dot product term in equation (1b) averaged to zero. This approximation will be adequate for the part of the neutrino effective potential generated by the electron current anywhere between the neutrino sphere and the shock, so long as there are no large-scale fluid motions. By contrast the neutrino distribution functions above the neutrino sphere are not isotropic, as the neutrinos are nearly freely streaming there, with the result that for $\nu\nu$-neutral current exchange the first two terms in equation (1b) tend to cancel and would completely cancel in the limit where the neutrinos move only on radial world lines at very large radius. The $\langle A^2\rangle$ and $\langle\phi^2\rangle$ terms are second-order weak interactions and are therefore completely negligible.

The effective mass contribution to a propagating ν_e from forward exchange scattering on the electron sea is

$$m_{\rm eff}^2 \approx 2E_{\nu_e}\langle\phi\rangle \approx (1.5184 \times 10^3\ {\rm eV}^2)(\rho_{10} Y_e)\left(\frac{E_{\nu_e}}{\rm MeV}\right) , \quad (2a)$$

where ρ_{10} is the density in units of 10^{10} g cm^{-3} and Y_e is the number of electrons per baryon. The analogous effective mass contribution from forward exchange scattering on the ν_e and $\bar{\nu}_e$ seas at radius R is roughly

$$\begin{aligned} m_{\rm eff}^2 &\approx 2E_\nu\langle\phi\rangle - 2\langle \boldsymbol{p}_{\nu_e} \cdot \boldsymbol{A}\rangle \\ &\approx 2E_\nu\langle\phi\rangle\left(\frac{R_{\nu s}}{R}\right)^2 \\ &\approx (1.5184 \times 10^3\ {\rm eV}^2)(n_{\nu_e} - n_{\bar{\nu}_e})\left(\frac{E_\nu}{\rm MeV}\right)\left(\frac{R_{\nu s}}{R}\right)^2 , \end{aligned} \quad (2b)$$

where $R_{\nu s}$ is the radius of the neutrino sphere and n_{ν_e} and $n_{\bar{\nu}_e}$ are the appropriate fluxes (or number densities) at radius R. Note that the scattering amplitude for $\nu_e\nu_e$-exchange is the negative of that for $\nu_e\bar{\nu}_e$-exchange (cf. Savage, Malaney, & Fuller 1991).

Equation (2b) shows that in most cases of interest we can neglect $\nu\nu$-exchange-scattering contributions to neutrino effective masses. If the core mass is of order 1.5 $M_\odot$, then a net electron-neutrino lepton number of $\approx 9 \times 10^{56}$ escapes in roughly ~ 10 s, so that the net effective neutrino density, $n_{\nu_e} - n_{\bar{\nu}_e}$, at a typical neutrino level crossing radius ($R \approx 160$ km for $E_\nu \approx 35$ MeV) is 1.2×10^{31} cm^{-3} which, with the geometric term in equation (2b), yields an effective number of scattering centers of 9×10^{29} cm^{-3}. This is to be compared with the ambient electron density at the same radius at, for example, 0.17 s pb (where pb stands for postbounce) in the R. Mayle & J. R. Wilson (1991, private communication) calculation where $n_e \approx 1.8 \times 10^{32}$. We can conclude that electrons dominate over neutrinos in effective mass contribution by a factor of about 200. Changing the outgoing flux of neutrinos by a factor of 10 either way by, for example, having the net lepton number escape in 1 s would not change this conclusion. Now there may be *local* increases in the net neutrino-lepton number flux of more than an order of magnitude or two, depending on the deleptonization time scale, so that we should be aware of $\nu\nu$-exchange-scattering effects. These effects become more important at radii close to the neutrino sphere.

Following the discussion in FMWS87 we can relate the instantaneous neutrino flavor eigenstates $|\nu_e(t)\rangle$ and $|\nu_x(t)\rangle$ at time t to the instantaneous mass eigenstates $|\nu_1(t)\rangle$ and $|\nu_2(t)\rangle$ by

$$\begin{aligned} |\nu_e(t)\rangle &= \cos\varphi|\nu_1(t)\rangle + \sin\varphi|\nu_2(t)\rangle , \\ |\nu_x(t)\rangle &= -\sin\varphi|\nu_1(t)\rangle + \cos\varphi|\nu_2(t)\rangle , \end{aligned} \quad (3)$$

where ν_x is either a ν_μ or a ν_τ and φ is the effective mixing angle in matter. In vacuum $\varphi = \theta$ (where θ is the vacuum mixing angle), and m_1 and m_2 are the vacuum mass eigenvalues. We define $\Delta \equiv m_1^2 - m_2^2$. A level crossing, or resonance, occurs when the matter contributions to the ν_e-ν_x effective mass difference, $\Delta_{\rm eff} = m_{\rm eff}^2$, are equal to $\Delta\cos\theta$. If, as we expect, the vacuum mixing angles are small and the vacuum masses obey the relation $m_{\nu_e} \ll m_{\nu_x}$, then $\Delta \approx m_{\nu_x}^2$. The instantaneous matter mixing angle, oscillation length, and mass eigenvalues are given in equations (9a–9d) in FMWS87. The level crossing condition implies that the resonant density is

$$\rho_{\rm res} \approx (2.108 \times 10^{10}\ {\rm g\ cm^{-3}})\left(\frac{0.5}{Y_e}\right)\left(\frac{\Delta}{1600\ {\rm eV}^2}\right)\left(\frac{\rm MeV}{E_\nu}\right) , \quad (4)$$

where E_ν is the energy of the neutrino and Y_e is the number of electrons per baryon. Note that this expression neglects $\nu\nu$-exchange-scattering contributions to the neutrino effective mass. If $\nu\nu$-exchange scattering is not negligible, then its contributions to neutrino effective mass have the sense of pushing the mass level crossing out to lower density in the supernova. The oscillation length at resonance is

$$L_{\rm res} = \frac{4\pi E_\nu}{\Delta\sin 2\theta} \approx \frac{0.1575\ {\rm cm}(E_\nu/{\rm MeV})}{(\Delta/1600\ {\rm eV}^2)\sin 2\theta} , \quad (5)$$

while the resonance "width" (Bethe 1986) is

$$\delta r = \left(\frac{1}{\rho}\frac{d\rho}{dr}\right)^{-1} \tan 2\theta . \quad (6)$$

A schematic picture of the neutrino emission and relative positions of the neutrino sphere, shock, and resonant region is given in Figure 1. This figure is based on a numerical result at 0.17 s pb in the Mayle & Wilson (1991) calculation. This computation includes a complete nuclear and electron equation of state, hydrodynamics, and a detailed treatment of neutrino transport (see Mayle 1990, and references therein). The neutrino sphere is near the edge of the hot proto–neutron star at about 47 km in radius for ν_e and $\bar{\nu}_e$, and at about 43 km for

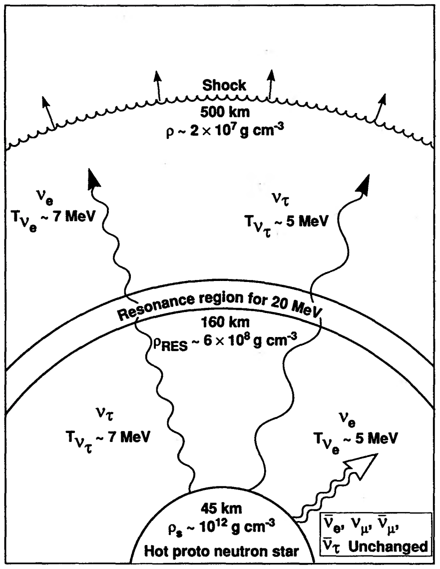

FIG. 1.—Schematic picture of the relative positions of neutrino sphere, shock, and resonant region in a numerical supernova model at a time 0.17 s pb. The level crossing position for a $\Delta = 1600$ eV^2 neutrino mass difference is shown for a neutrino energy of $E_\nu = 20$ MeV. The density at the level crossing, or resonance density $\rho_{\rm RES}$, for these parameters is also given.

$\nu_\tau(\bar{\nu}_\tau)$ and $\nu_\mu(\bar{\nu}_\mu)$. Because $\nu_\mu(\bar{\nu}_\mu)$ and $\nu_\tau(\bar{\nu}_\tau)$ have only neutral current interactions they decouple deeper in the core than do the $\nu_e(\bar{\nu}_e)$ which have charged as well as neutral current interactions. The position of each neutrino sphere corresponds, roughly, to a density of about $\rho Y_e \approx 10^{12}$ g cm^{-3}, while the position of the shock at this epoch corresponds to a density $\rho Y_e \approx 2 \times 10^7$ g cm^{-3}.

From Figure 1 and equation (4) it can be seen that a level crossing of ν_e with ν_x occurring above the neutrino sphere but below the stalled shock would require a vacuum mass difference corresponding to $\Delta \approx 10^2$–10^4 eV^2, which, if m_{ν_e} is negligible, corresponds to $m_{\nu_x} \approx 10$–100 eV. If we define Ω_ν to be the fraction of the closure density of the universe contributed by neutrinos, then the mass m_{ν_x} of the neutrino must be

$$m_{\nu_x} \approx (96\ {\rm eV})\left(\frac{H_0}{100\ {\rm km\ s^{-1}\ Mpc^{-1}}}\right)^2\left(\frac{2.7\ {\rm K}}{T_\gamma}\right)^3 \Omega_\nu\,, \quad (7)$$

where H_0 is the Hubble constant and T_γ is the present cosmic microwave background temperature. Reasonable ranges of H_0 and Ω_ν then yield a cosmologically interesting mass range for m_{ν_x} which is coincident with the level crossing range discussed above for supernovae. The popularity of a neutrino as a dark matter candidate particle has waxed and waned with time, principally due to the difficulties in understanding how galaxy formation could occur in a hot-dark-matter top-down scheme (cf. Cowsik & McClelland 1972; Gunn et al. 1978; Tremaine & Gunn 1979; Bond, Efstathiou, & Silk 1980; Blumenthal et al. 1984). However, we feel that it is unwise to rule out the possibility of a closure mass neutrino on the basis of the observed structure distribution in the universe until the galaxy formation process is better understood. Ideas on galaxy formation are far from being definitive, so at this point constraints derived from these considerations neither encourage or discourage us.

The probability that a ν_e transforms to a ν_x in the resonance region depends on a comparison of $L_{\rm res}$ and δr. In the adiabatic limit the propagating neutrino sees such a gradual change in density or, equivalently, $m^2_{\rm eff}$ that $|\nu_e\rangle$ is completely rotated into $|\nu_x\rangle$ and vice versa. Haxton (1987) (see also Parke & Walker 1986) has given a beautiful derivation of this result which we employ here. The Landau-Zener jump probability is $P_{\rm LZ}(E_{\nu_e})$, where

$$P_{\rm LZ}(E_{\nu_e}) \approx \exp\left(-\frac{\pi^2}{2}\frac{\delta r}{L_{\rm res}}\right), \quad (8a)$$

so that the adiabatic condition is $L_{\rm res} \ll \delta r$. The probability that a ν_e of energy E_{ν_e} transforms to a ν_x of the same energy and vice versa is

$$P(E_{\nu_e}) \approx 1 - P_{\rm LZ}(E_{\nu_x})\,. \quad (8b)$$

In practice then $\nu_e \rightleftharpoons \nu_x$ transformation is important whenever the vacuum mixing angle satisfies

$$\frac{\sin^2 2\theta}{\cos 2\theta} \geq \frac{8}{\pi}\left(\frac{1}{\rho}\frac{d\rho}{dr}\right)\frac{E_\nu}{\Delta}\,, \quad (9a)$$

where $[(1/\rho)(d\rho/dr)]$ is to be evaluated in the supernova at the level crossing point. If $\theta \ll 1$ then this condition becomes

$$\theta \geq \left(\frac{2E_\nu}{\pi\Delta}\right)^{1/2}\left(\frac{1}{\rho}\frac{d\rho}{dr}\right)^{1/2}. \quad (9b)$$

For typical conditions ($\Delta \approx 1600$ eV^2), the density scale height at resonance is $[(1/\rho)(d\rho/dr)]^{-1} \approx 5 \times 10^6$ cm so that the condition in equation (9b) becomes

$$\theta \geq 10^{-4}\left(\frac{E_\nu}{35\ {\rm MeV}}\right)^{1/2}. \quad (9c)$$

Representative neutrino energy spectra for ν_e and ν_τ at their respective neutrino spheres are shown in Figure 2. These distribution functions are taken from the Mayle & Wilson (1991) calculation at a time 0.1 s pb, coinciding with the beginning of the reheating epoch. The energy spectra for $\bar{\nu}_\tau$, ν_μ, and $\bar{\nu}_\mu$ are essentially identical to that for ν_τ. The spectrum for $\bar{\nu}_e$ is similar to that for ν_e (but about 25% hotter) at this time, although as discussed above, there occasionally may be significant contribution to the net ν_e flux from neutronization neutrinos which would cause the ν_e and $\bar{\nu}_e$ spectra to differ. In the absence of

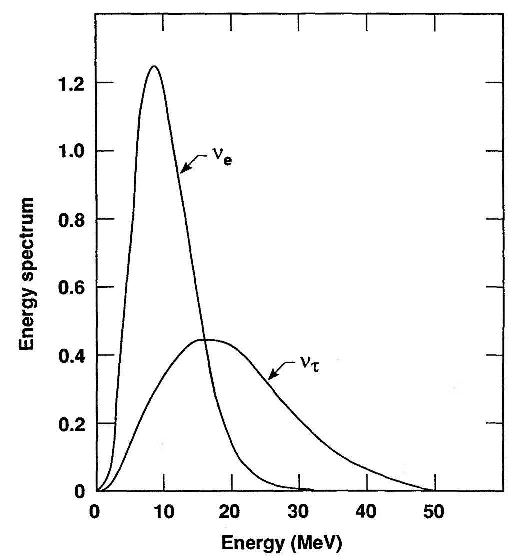

FIG. 2.—Representative neutrino energy spectra for ν_e and ν_τ near their neutrino spheres. From a numerical supernova model at 0.1 s pb.

TABLE 1

NEUTRINO ENERGY GROUPS

Group	Central Energy (MeV)
1	2.000
2	2.828
3	3.999
4	5.654
5	7.995
6	11.31
7	15.99
8	22.60
9	31.96
10	45.19
11	63.90
12	90.36
13	127.8
14	180.7
15	255.5
16	361.2

significant neutrino flavor transformations it would be nearly correct to think of the neutrino distribution functions at radii larger than the neutrino-sphere radius to correspond to those of the neutrino sphere at the appropriately retarded time.

If neutrino mass level crossings occur in the region between the neutrino sphere and the radius in question, then neutrino flavor transformations may modify the above prescription for computing the neutrino distribution functions. In this case computation of the distribution functions would require that we follow the phases of neutrinos at each energy through the resonance region. To facilitate this calculation, as well as transport calculations, we represent the neutrino distribution function for each species by a radius and time-dependent occupation number for each of 16 energy groups or bins. Table 1 presents the neutrino energies in MeV corresponding to each energy group used in a typical calculation. Each neutrino energy group will have a resonance (level crossing) at a density given by equation (4). We compute the neutrino transformation probability at these resonances in the manner prescribed in equations (5)–(9). All of these calculations are carried out implicitly within the Mayle & Wilson (1991) supernova code along with neutrino transport and hydrodynamics. This method for calculating neutrino transformation rates will accurately serve our purposes so long as the neutrino vacuum mixing angles are small enough to ensure that the resonance widths are small compared to the reheating region length scales we are interested in (typically a few zone widths).

We implement the neutrino transformation calculations in our supernova code as follows. It will turn out that in the reheating epoch the resonance densities for the neutrino masses of interest are at most of order 10^{10} g cm^{-3} and, hence, correspond to radii which lie well outside the neutrino sphere, as depicted in Figure 1. The neutrinos are assumed to be outwardly free streaming through the resonance regions. For each neutrino energy group in the calculated supernova model we go to the zone whose density is closest to the resonance density as computed from equation (4). There we implement the changes in occupation numbers of each neutrino energy group as follows: we convert a fraction,

$$f = 1 - P_{LZ}(E_\nu) , \tag{10}$$

of all entering ν_x into ν_e. Likewise, a fraction f of the entering ν_e are converted to ν_x. The Landau-Zener transformation probability is given by equation (8a), where the logarithmic derivative of density in equation (6) is computed from the data in adjacent zones of the numerical supernova model. This neutrino transformation procedure involving only a single zone is strictly quantum mechanically correct only if the resonance width is less than the zone size, although little error would be introduced in our reheating calculations were this criterion to be relaxed somewhat.

We can gauge the dependence of the completeness of neutrino flavor transformations on vacuum mixing angles and masses by examining the Mayle & Wilson (1991) calculation at times corresponding to the supernova configuration shown in Figure 1. Table 2 presents the radius of the level crossing, the density scale height at resonance (resonance width divided by tan 2θ), and the ratio of the resonance width to the oscillation length with the vacuum mixing-angle dependence scaled out, all for each neutrino energy in a supernova model at 0.17 s pb. These quantities are given for two $\nu_e - \nu_x$ mass differences: $\Delta = 225$ eV^2 and $\Delta = 3600$ eV^2. Table 3 presents the same

TABLE 2

LEVEL CROSSING PARAMETERS

($t$pb = 0.177 s)

Energy Group	Radius of Level Crossing (in 10^6 cm)		$\frac{\delta r}{\tan 2\theta} = \left(\frac{1}{\rho}\frac{d\rho}{dr}\right)^{-1}$ (in 10^6 cm)		$\frac{\delta r}{L_{res}}\left(\frac{\cos 2\theta}{\sin^2 2\theta}\right) \times 10^{-6}$	
	$\Delta = 225$ eV^2	$\Delta = 3600$ eV^2	$\Delta = 225$ eV^2	$\Delta = 3600$ eV^2	$\Delta = 225$ eV^2	$\Delta = 3600$ eV^2
1	12.5	6.5	4.083	1.013	1.82	7.2
2	12.7	7.0	3.06	1.11	0.965	5.62
3	15.6	6.9	6.84	0.787	1.53	2.81
4	17.8	7.8	6.94	1.509	1.09	3.81
5	20.0	8.4	7.1	1.89	0.794	3.38
6	23.5	8.9	10.8	2.412	0.855	3.05
7	27.8	9.7	15.7	2.91	0.875	2.6
8	32.5	11.1	14.9	3.07	0.59	1.94
9	39.0	12.2	19.87	4.33	0.56	1.93
10	46.0	13.7	11.7	4.37	0.23	1.38
11	48.0	15.5	8.27	4.66	0.11	1.04
12	51.0	17.8	7.6	6.95	0.075	1.1
13	54.0	19.0	5.38	7.11	0.0376	0.795
14	55.0	23.8	3.8	10.85	0.019	0.86
15	56.6	27.0	12.8	11.07	0.045	0.62
16	62.0	31.0	9.03	0.149	0.023	0.58

TABLE 3

LEVEL CROSSING PARAMETERS

($t_{pb} \simeq 0.638$ s)

Energy Group	Radius of Level Crossing (in 10^6 cm)		$\frac{\delta r}{\tan 2\theta} = \left(\frac{1}{\rho}\frac{d\rho}{dr}\right)^{-1}$ (in 10^6 cm)		$\frac{\delta r}{L_{res}}\left(\frac{\cos 2\theta}{\sin^2 2\theta}\right) \times 10^{-6}$	
	$\Delta = 225$ eV^2	$\Delta = 3600$ eV^2	$\Delta = 225$ eV^2	$\Delta = 3600$ eV^2	$\Delta = 225$ eV^2	$\Delta = 3600$ eV^2
1.........	2.7	2.3	0.204	0.04906	0.091	0.35
2.........	2.8	2.35	0.261	0.0815	0.0825	0.411
3.........	2.9	2.36	0.271	0.0576	0.061	0.206
4.........	3.0	2.37	0.419	0.1324	0.066	0.334
5.........	3.5	2.4	0.488	0.0936	0.0544	0.167
6.........	3.6	2.44	0.345	0.252	0.0272	0.3183
7.........	3.7	2.52	0.453	0.320	0.0253	0.286
8.........	3.9	2.62	0.653	0.289	0.0258	0.183
9.........	4.0	2.7	0.462	0.2041	0.01289	0.09125
10.........	4.4	2.8	0.7354	0.262	0.01453	0.0828
11.........	6.2	2.9	0.983	0.272	0.013	0.06075
12.........	6.4	3.0	0.696	0.4194	0.00687	0.0663
13.........	6.7	3.44	1.44	0.488	0.0101	0.0546
14.........	7.2	3.6	1.021	0.3453	0.005047	0.0273
15.........	9.8	3.75	3.70	0.4537	0.0129	0.02535
16.........	11.0	3.9	2.86	0.653	0.00707	0.02584

quantities for this supernova model but for an epoch near the end of the first reheating phase at 0.64 s pb. The radius of the shock at the epoch corresponding to Table 2 is ≈ 500 km, whereas, the radius is ≈ 4000 km for the time corresponding to Table 3.

Several interesting trends can be gleaned from Tables 2 and 3. First we can see that the resonance width for a ν-energy near the peak of the ν_τ distribution in Figure 2 will be small (< 1 km) so long as the vacuum mixing angle is $\theta < 2 \times 10^{-2}$ early in the reheating process and $\theta < 0.2$ later in this epoch. This justifies our procedure for computing the neutrino transformation rates as outlined above. We can also conclude from these tables that the estimate of the minimum vacuum mixing angle required for adiabaticity in ν-transformation as given in equation (9c) is roughly correct. Note, however, that the highest energy neutrinos will have a lower transformation probability (be less "adiabatic") than lower energy ones at a given vacuum mixing angle. This is to be expected as the oscillation length at resonance increases with neutrino energy, while the supernova models yield a density scale height which varies only slowly with radius and time, at least in the region where reheating is most important. We will examine the effect of our neutrino oscillation calculations on the shock revitalization process after we have discussed how charged-current neutrino interactions can cause energy deposition and a significant heating of the matter behind the shock.

3. CONVENTIONAL SUPERNOVA SHOCK REHEATING BY NEUTRINOS

In this section we will examine how neutrino energy emitted from a postbounce hot-proto–neutron-star core might be coupled into the internal energy of the material behind the shock. We will assume that the supernova shock has moved out 300 to 500 km from the core and stalled. There are two subsequent shock reheating epochs. First, commencing at roughly 0.1 s pb, and proceeding out to 0.6 s pb, ν_e and $\bar{\nu}_e$ can be captured on neutrons and protons,

$$\nu_e + n \rightarrow p + e^- , \quad (11a)$$

$$\bar{\nu}_e + p \rightarrow n + e^+ , \quad (11b)$$

to heat material behind the shock, locally increasing the pressure and, thus, the Mach number of the shock. This reheating process was considered in some detail in BW85, and we will follow their notation here. Bruenn (1991) has also studied this process with results which are sometimes at odds with the findings of Mayle & Wilson (1991) and BW85. Although the physics of shock revival effected by neutrino interactions is an extremely complex problem which may not be resolved for some time, our study is designed to ascertain the added effects of putative neutrino oscillations. The second possible reheating epoch comes at later times (> 0.5 s pb), after hot bubble formation, and is due to neutrino pair annihilation, neutrino-electron scattering, as well as neutrino charged-current captures feeding energy into the hot plasma above the neutrino sphere (Colgate 1991; Goodman, Dar, & Nussinov 1987).

From Figure 1, which corresponds roughly to the onset of the first reheating phase, we see that the stalled shock has a radius of ≈ 500 km, where the density is $\rho Y_e \approx 2 \times 10^7$ g cm^{-3}. We note that the position of the shock at this time is quite sensitive to the initial shock energy, the amount of material in the outer core which must be photodisintegrated, and, for example, the magnitudes of the various neutrino fluxes. Efficiency of reheating will increase with increasing shock radius at this epoch.

The various neutrino energy spectra at this epoch are roughly blackbody (Fermi-Dirac, zero chemical potential), but fall below this at high neutrino energy. The characteristic neutrino-sphere spectral temperatures $T_{\nu_e(\bar{\nu}_e)}$ and $T_{\nu_\mu(\nu_\tau)}$ vary with individual calculations but are roughly $T_{\nu_e} \approx T_{\bar{\nu}_e} \approx 5$ MeV and $T_{\nu_\mu(\nu_\tau)} \approx 7$ MeV at ≈ 0.2 s pb. Actually, numerical calculations show $T_{\bar{\nu}_e}$ to be about an MeV higher than T_{ν_e} at this epoch. For the simple analytic estimates we wish to make here it suffices to take these temperatures as being the same.

The neutrino-energy–dependent cross sections for processes (11a) and (11b) are

$$\sigma(E_\nu) = \langle G \rangle \frac{\ln 2}{\langle ft \rangle}\left[\frac{2\pi^2(\hbar c)^3}{c}\right](m_e c^2)^{-5}(Q_n + E_\nu)^2 , \quad (12a)$$

$$\approx (9.23 \times 10^{-44}\ \text{cm}^2)\langle G \rangle \left(\frac{Q_n + E_\nu}{\text{MeV}}\right)^2 , \quad (12b)$$

where $\hbar$ is Plancks' constant, c is the speed of light, $\log_{10}\langle ft \rangle \approx 3.035$ is the effective ft-value for free nucleons

(Fuller, Fowler, & Newman 1985), and Q_n is the difference between parent and daughter *bare nucleon* masses,

$$Q_n = \pm(m_n - m_p) \approx \pm 1.293 \text{ MeV} , \quad (12c)$$

where the plus is for process (11a) and the minus for process (11b). In equation (12) the average correction for Coulomb waves is $\langle G \rangle$, so that $\langle G \rangle = 1$ for process (11b) and is given in Fuller et al. (1980) as $\langle G \rangle = G_+ \approx 1$ for process (11a). The energy-dependent neutrino opacity, or mass absorption coefficient, corresponding to $\sigma(E_\nu)$ is

$$K_i = N_A Y_i \langle \sigma(E_\nu) \rangle , \quad (13a)$$

where N_A is Avogadro's number, i the runs over neutron or proton target, and Y_i the appropriate nucleon number per baryon. Then Y_n, Y_p, and Y_N are the free neutron, free proton, and total free nucleon number per baryon, respectively. The bracket denotes an average over the actual neutrino energy spectrum. BW85 found that for ν_e and $\bar{\nu}_e$ the mass absorption coefficient is roughly

$$K_i(T_\nu) \approx (3.8 \times 10^{-19} \text{ cm}^2 \text{ g}^{-1}) Y_i T_\nu^2 , \quad (13b)$$

which partially takes account of the paucity of high-energy neutrinos relative to a blackbody.

The net specific heating rate (ergs g^{-1} s^{-1}) for material immediately behind the shock is then

$$\dot{E}_{\rm BW85} \approx (4\pi R_m^2)^{-1}[K_n(T_{\nu_e})L_{\nu_e} + K_p(T_{\bar{\nu}_e})L_{\bar{\nu}_e}] - 4\pi j(T_m) , \quad (14a)$$

where L_{ν_e} and $L_{\bar{\nu}_e}$ are the total ν_e and $\bar{\nu}_e$ luminosities, R_m and T_m refer to the radius and local temperature, respectively, of a matter element behind the shock, and $j(T_m)$ is the neutrino emissivity per steradian of material at temperature T_m. Since we take $T_{\nu_e} \approx T_{\bar{\nu}_e}$, and both temperatures are usually large enough that $K_n \approx K_p$ and $Y_N = Y_n + Y_p$ (so we can define $K = K_n + K_p$), and $L_{\nu_e} \approx L_{\bar{\nu}_e}$, we can approximate the result in equation (14a) as (BW85)

$$\dot{E}_{\rm BW85} \approx K(T_{\nu_e})\left[\frac{L_\nu}{4\pi R_m^2} - \left(\frac{T_m}{T_\nu}\right)^2 acT_m^4\right] , \quad (14b)$$

where we have used detailed balance and the steady state condition to rewrite the loss term as $4\pi j(T_m) = K(T_m)cu_\nu$, with u_ν the *equivalent* thermal energy density in neutrinos at temperature $T = T_m$,

$$u_\nu(T) = aT^4 = \frac{7}{8} g_f\left(\frac{\pi^2}{30}\right)T^4 , \quad (14c)$$

where g_f is the neutrino statistical weight, with $g_f = 1$ for each neutrino species. The neutrino energy loss term from $e^\pm$ capture in equation (14) is negligible when $T_m \ll T_{\nu_e}$. This condition will not be met at the beginning of reheating but is not a bad approximation once we reach times after 0.4 s pb. The matter temperature, T_m, at a given position *initially* is determined by the gravitational potential there. This is because the material upstream of the shock eventually flows through the shock with nearly the free-fall velocity, whereupon the material is decelerated and its kinetic energy is converted to internal energy. The initial postshock-material temperature distribution is then set by the mass distribution in the precollapse star. Thus the heating process can be aided primarily by increasing T_ν and L_ν.

If there are neutrino oscillations which effect ν_e but not $\bar{\nu}_e$ (as is expected for a neutrino mass hierarchy where $m_{\nu_\tau} > m_{\nu_\mu} > m_{\nu_e}$), then L_{ν_e} and $L_{\bar{\nu}_e}$, as well as $K_n(T_{\nu_e})$ and $K_p(T_{\bar{\nu}_e})$, may differ appreciably from their nonoscillation values. This would result in considerable alteration in the neutrino heating rate.

4. SHOCK REHEATING WITH NEUTRINO OSCILLATIONS

Having discussed neutrino energy deposition and shock revival in conventional delayed-mechanism supernova models, we will now examine how the picture might be expected to change if matter-enhanced transformations among neutrino flavors are important in the reheating region. We will first make a simple analytic estimate along the lines of that given in the last section and then present the results of detailed numerical calculations performed with a coupled neutrino-oscillation/supernova-hydrodynamics code.

We can get a rough idea of the effect of neutrino oscillations on shock revival without resorting to numerical calculations by approximating the neutrino distributions as nearly blackbody in character and by taking the $\nu_e \rightleftharpoons \nu_x$ transformation at resonance as complete. The BW85 technique can then be employed to contrast the net-energy deposition rates with and without neutrino oscillations. This set of approximations is a fair estimate for reheating calculations since the most effective range of neutrino energies for energy deposition is of order $3T_\nu \sim 10$–35 MeV. Neutrinos more energetic than this may have a smaller transformation probability (eqs. [8]–[10]), and the neutrino energy spectra at high energy deviate more significantly from blackbody distributions. If $\theta > 10^{-4}$ then over the energy range of interest complete neutrino flavor transformation is a fair approximation. Neutrinos with energies below this range have small cross sections and do not contribute much to heating. Neutrinos with very high energies are not important for heating since their fluxes are low. Roughly then the temperature which characterizes the ν_e energy distribution in the relevant energy range is now T_{ν_x}, the neutrino-sphere temperature for the ν_x, while the temperature which characterizes the ν_x distribution is now to T_{ν_e}. Of course, as pointed out above, $T_{\nu_x} > T_{\nu_e}$.

Concentrating on the heating term alone in equation (14) we can write the heating rate, $\dot{E}_+$, as

$$\dot{E}_+ \approx \frac{L_{\bar{\nu}_e}}{4\pi R_m^2} K(T_{\bar{\nu}_e})\left[Y_p + Y_n \frac{L_{\nu_e}}{L_{\bar{\nu}_e}} \frac{K(T_{\nu_e})}{K(T_{\bar{\nu}_e})}\right] , \quad (15a)$$

where we neglect the difference in Q_n for neutrons and protons in equation (12). Now if the ν_e and ν_x are completely transformed, one into another in the relevant energy range, then we can approximate $T_{\nu_e} \to T_{\nu_x}$ so that

$$\frac{\dot{E}_+}{(\dot{E}_+)_{\rm BW85}} \approx \frac{1}{Y_N}\left[Y_p + Y_n\left(\frac{T_{\nu_x}}{T_{\bar{\nu}_e}}\right)^6\left(\frac{r_{\nu_x}}{r_{\nu_e}}\right)^2\right] , \quad (15b)$$

where the heating rate without oscillations as estimated in BW85 is

$$(\dot{E}_+)_{\rm BW85} \approx \frac{L_{\bar{\nu}_e}}{4\pi R_m^2} Y_N K(T_{\bar{\nu}_e}) , \quad (15c)$$

and where r_{ν_e} and r_{ν_x} are the neutrino-sphere radii for ν_e and ν_x, respectively. For the Mayle & Wilson (1991) model near the beginning of the reheating epoch at 0.17 s pb we find that $r_{\nu_e} \approx 47$ km, while $r_{\nu_x} \approx 43$ km. Typical temperatures for this epoch are $T_{\bar{\nu}_e} \approx 5$ MeV and $T_{\nu_x} \approx 7$ MeV, while the nucleon mass fractions are roughly $Y_N \approx 1$, $Y_p \approx \frac{1}{3}$, and $Y_n \approx \frac{2}{3}$, so that equation (15) would predict approximately a factor of 4 increase in the heating rate with oscillations over the BW85 result given by equation (15c). However, numerical calculations show that after a few tens of milliseconds the luminosities of all the neutrino species are approximately equal. Setting $L_{\nu_x} \approx L_{\nu_e}$ in equation (15a) then yields a more realistic estimate for the ratio of the heating rates with and without oscillations:

$$\frac{\dot{E}_+}{(\dot{E}_+)_{\rm BW85}} \approx \frac{1}{Y_N}\left[Y_p + Y_n\left(\frac{T_{\nu_x}}{T_{\bar{\nu}_e}}\right)^2\right]. \qquad (15d)$$

This expression would predict about a factor of 2 increase in heating rate over the BW85 result. This is clearly a very idealized treatment. The neutrino heating rate, along with energy loss and dissipation terms, would have to be integrated over time to give the actual shock energy and total supernova explosion energy. Nonlinear effects, especially the material-neutrino-loss term neglected above, would be expected to be very important.

We perform our numerical neutrino-oscillation/supernova-hydrodynamics calculations with a variant of the Mayle & Wilson (1991) code for a model with an initial Fe-core mass of about 1.43 $M_\odot$. A prescription for computing the dynamic evolution of the neutrino distribution functions is included which is essentially similar to that described in § 2. Since, as discussed above, the neutrino energies which will be most important for the reheating process are of order 10–40 MeV, we have used a somewhat finer energy resolution between neutrino energy groups than that in Table 1. Of course we do not use Fermi-Dirac blackbody distributions, but rather the actual neutrino energy spectra. Table 4 presents the central energies for the neutrino energy groups (in MeV) in a representative calculation and gives the spectral occupation numbers (equivalent to a measure of flux) for ν_e and one of either ν_μ, $\bar{\nu}_\mu$, ν_τ, or $\bar{\nu}_\tau$ (in this calculation they are all the same, so that we label the column as F_{ν_τ}) at times 0.10, 0.20, 0.30, and 0.45 s pb.

Our numerical calculations are all done for a fixed vacuum mixing angle $\theta = 10^{-3}$ between ν_e and ν_x. Of course ν_x could be either ν_μ or ν_τ, but in our treatment the results are independent of which one is chosen. However, the expected neutrino signal in some supernova detectors may well depend on which of these neutrinos transforms in the reheating region, as the other neutrino may well have a level crossing with the ν_e which are located further out in the star. The numerical oscillation calculations with our selected vacuum mixing angle yield nearly adiabatic neutrino flavor transformations in the reheating region. Our results, at least for neutrino heating, should be similar for any vacuum mixing angle $10^{-4} < \theta < 2 \times 10^{-2}$.

With the vacuum mixing angle fixed, each neutrino energy group has a resonance density determined by Δ as in equation

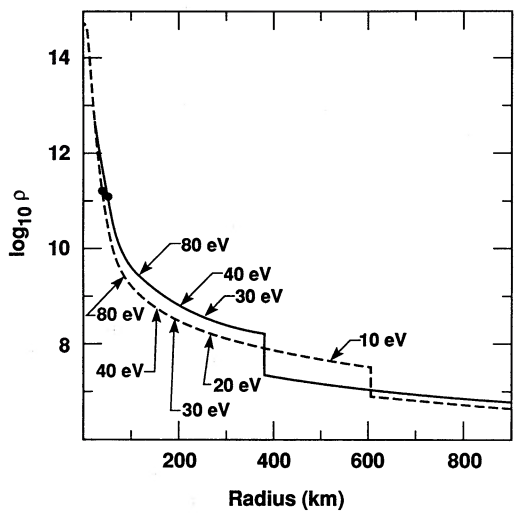

FIG. 3.—The run of density (ρ in g cm^{-3}) with radius (in km) for a numerical supernova model at a time 0.15 s pb (*solid curve*), and at 0.2 s pb (*dashed curve*). The position of the neutrino sphere is represented as a filled circle at each time. Marked on each curve are the neutrino mass level crossing, or resonance, densities for various $\Delta^{1/2}$ measured in eV.

(4). If, as outlined above, $m_{\nu_x} \gg m_{\nu_e}$, then $\Delta^{1/2}$ is a measure of the ν_x vacuum mass. Figures 3 and 4 show the run of density with radius for several numerical supernova models. Marked on each figure are the level crossing, or resonance, densities for various $\Delta^{1/2}$ measured in eV. The positions of the neutrino spheres are shown as filled circles. In Figure 3 the solid curve corresponds to a time 0.15 s pb, while the dashed curve represents the same model at a time 0.2 s pb. The discontinuity in density corresponding to the shock is clearly visible in these models, and we see that the shock moves outward in radius by a small amount during the time elapsed between models. Figure 4 corresponds to the same supernova model but now at an epoch 0.45 s pb, by which time the shock has moved out to a radius of about 3000 km.

In order for matter-enhanced neutrino flavor transformations to deliver a higher ν_e flux with a harder spectrum to the material behind the shock we obviously must have the neutrino mass level crossing occur at a radius well inside that of the shock. This favors a higher $\Delta^{1/2}$, or m_{ν_x}, within the cosmologically interesting mass range discussed above, at least as far as reheating efficiency goes. We have performed detailed reheating calculations as outlined above for $\Delta^{1/2} = 40$ and

TABLE 4

NEUTRINO SPECTRA

	$t_{\rm pb} = 0.10$ s		$t_{\rm pb} = 0.20$ s		$t_{\rm pb} = 0.30$ s		$t_{\rm pb} = 0.45$ s	
E_ν (MeV)	F_{ν_e}	F_{ν_τ}	F_{ν_e}	F_{ν_τ}	F_{ν_e}	F_{ν_τ}	F_{ν_e}	F_{ν_τ}
4.00..........	4.39	0.795	1.38	0.28	0.733	0.18	0.335	0.07
5.04..........	6.82	1.26	2.31	0.47	1.23	0.30	0.567	0.12
6.35..........	9.80	1.91	3.63	0.76	1.95	0.49	0.900	0.20
8.00..........	12.14	2.66	5.111	1.13	2.78	0.75	1.276	0.31
10.08..........	12.12	3.43	6.15	1.56	3.39	1.06	1.522	0.45
12.70..........	9.03	4.06	5.95	1.98	3.36	1.39	1.447	0.61
16.00..........	4.59	4.40	4.38	2.29	2.59	1.64	1.066	0.75
20.16..........	1.44	4.13	2.28	2.33	1.50	1.71	0.590	0.81
25.40..........	0.231	3.08	0.750	1.97	0.602	1.50	0.230	0.72
32.00..........	0.016	1.61	0.134	1.27	0.153	1.02	0.057	0.50
40.32..........	0.00036	0.48	0.0105	0.54	0.021	0.48	0.0080	0.25
50.80..........	...	0.06	...	0.13	0.094	0.13	0.00055	0.07
64.00..........	...	0.002	...	0.01	...	0.02	...	0.01

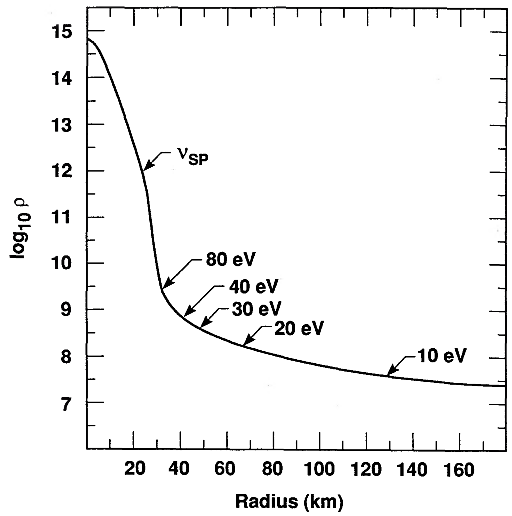

FIG. 4.—Same as Fig. 3, but now for the supernova model at a time 0.45 s pb.

90 eV and find in each case that the shock energy at very late times, and ultimately the supernova explosion energy, is increased by about 60% over a similar calculation without neutrino oscillations. This represents a significant increase in supernova energy, but is not as large as the factor of 2 increase predicted in the naive analytic estimate given above.

The reason for this discrepancy is clear: extra heating of the material behind the shock, due to a harder ν_e spectrum with a higher flux, would increase the temperature of the material. An increase in temperature would translate into higher thermal neutrino energy loss rates which would, in turn, reduce the net neutrino heating rate. This source of "negative feedback" in the heating process reduces the attainable neutrino heating efficiency.

5. CONSTRAINTS ON NEUTRINO OSCILLATIONS IN SUPERNOVAE

The detection of neutrinos from Supernova 1987A (SN 1987A) and the existence of large-scale accelerator and reactor neutrino oscillation experiments raise the question of whether the oscillation parameter space discussed here ($100 \text{ eV}^2 \leq \Delta \leq 10^4 \text{ eV}^2$ and $10^{-4} \leq \theta \leq 10^{-2}$) in connection with supernovae could be independently constrained. In this section we argue that the existing data are not sufficient to address the parameter space of interest, although future experiments and supernova neutrino detectors may be able to make a definitive statement on this subject.

Previous studies have attempted to constrain this parameter space by comparison of predicted neutrino energy spectra from supernova models which include neutrino oscillations with the 19 neutrino-induced events in the IMB and Kamiokande detectors for SN 1987A (cf. Arafune et al. 1987; Lagage et al. 1987; Kuo & Pantaleone 1988; Nötzold 1987). All of these studies assume, however, that the first two events in Kamiokande are ν_e-e scattering events and that the ν_e involved come from a distinct "neutronization pulse." These are both questionable assumptions (cf. Rosen 1988 for a discussion of the probability that the first two Kamiokande events are due to scattering; see also Burrows 1989 and Mayle 1990 for a general discussion of neutrinos from SN 1987A), so that any conclusions drawn are at best model dependent and therefore cannot exclude the parameter space of interest for neutrino-oscillation-enhanced reheating.

These assumptions are crucial to any inference of neutrino flavor transformation in SN 1987A for the following reasons. First, the signals in the IMB and Kamiokande detectors primarily result from $\bar{\nu}_e$-charged-current captures on protons, so that any ν_e-induced events must come from scattering. The indication that the first two events in Kamiokande are consistent with being forward peaked is taken as evidence that they are scattering events. If so, it is likely that they are ν_e-induced, since ν_e have a larger scattering cross section than do ν_μ or ν_τ, which interact through neutral current processes only. Rosen (1988) has argued, however, that with the small number of events involved here it is difficult to draw these conclusions with a high level of confidence. Note, however, that the second Kamiokande event is scattered at 40°. Thus, it may be argued that, at face value, this event is a poor candidate for a scattering event.

Now if these are indeed ν_e-e scattering events, then further interpretation would require that we know what stage of the proto–neutron-star cooling process the scattered neutrinos came from. The "neutronization pulse" occurs when the shock passes through the neutrino sphere. This pulse, while it has a large luminosity ($\lesssim 10^{54}$ ergs s^{-1}), persists only for 2–3 ms, which is a characteristic dynamical time at the neutrino sphere. The "neutronization pulse," preceding the thermal cooling phase, consists primarily of ν_e with mean energies less than about 10 MeV. If neutrino oscillations converted ν_e to ν_μ or ν_τ, then there would be a reduced scattering event rate observed in the detector. Observations of significant numbers of scattering events would then allow one to draw the conclusion that neutrino flavor transformation did not occur. Note that it is important for this argument that the scattered ν_e come from the "neutronization pulse," because if they come from the thermal pulse, where there are comparable fluxes of ν_μ and ν_τ, then the ν_e transformed to ν_x would just be replaced by the concomitant process of ν_x transforming to ν_e! The result would be no dimunition of ν_e-e scattering rate. Mayle (1990) has argued that it is unlikely that the first two Kamiokande events are from the neutronization pulse as they occur 0.1 s apart, and the duration of the "neutronization pulse" is considerably shorter (2–3 ms) as outlined above. Obviously much work needs to be done before any firm conclusions could be drawn regarding oscillation parameters or the complicated neutrino cooling history.

Kielczewska (1990) has pointed out another possible, more promising, constraint on oscillation parameters from observations of supernova neutrinos. This constraint is based on the backward-peaked nature of the ν_e-charged-current capture on ^{16}O (Haxton 1987, 1988) in water detectors like Kamiokande or IMB. Although the effective threshold for this capture process is high ($E_{\nu_e} > 30$ MeV to access substantial weak strength), the cross section rises very rapidly with energy due to the opening of extra forbidden-strength capture channels at high daughter-nucleus excitation energy (Haxton 1987, 1988; see also the discussion of angular distributions in neutrino capture in Fuller & Meyer 1991). This process may then become important for high-energy ν_e and thus may facilitate a useful constraint for the high end of the mixing-angle range in our oscillation parameter space.

If $\nu_e \rightleftharpoons \nu_x$ transformations are significant, especially for high-energy neutrinos, then we might expect an enhanced number of

$^{16}O(\nu_e, e^-)^{16}F$–induced backward-peaked events in a large water detector. Preliminary calculations in Qian, Fuller, & Meyer (1991) indicate, however that for SN 1987A we would expect only of order one of these events in, for example, Kamiokande if we assume that there is *complete* conversion of ν_x to ν_e. Given the poor statistics in the SN 1987A neutrino detections we cannot say with certainty whether or not this is consistent with the data. On the other hand these considerations are tantalizingly close to yielding an important bound. Perhaps the large water detectors which are being proposed for the future should be optimized to look for just such signals, as has been suggested by Kielczewska (1990).

We note from Tables 2 and 3 that for mixing angles $\theta \approx 10^{-4}$ the neutrinos which are important for reheating ($E_\nu \approx$ 10–40 MeV) might be nearly adiabatically transformed, while those which would make the greatest contribution to the ^{16}O-capture-induced events ($E_\nu > 40$ MeV) would not be transformed. Unfortunately this might provide a loophole, such that the absence of backward-peaked events could not then eliminate the possibility of ν-transformation in the important reheating energy range.

Minakata & Nunokawa (1990) have explored ways in which the expected time and energy signature of a supernova neutrino burst might be compared with the observed response in water detectors to yield constraints on neutrino oscillations. We note, however, that useful constraints would require very good detector event statistics and would therefore demand a much larger detector than currently exists.

Finally, we point out that although laboratory oscillation experiments involving $\nu_\tau \rightleftharpoons \nu_\mu$ conversion can probe the range of Δ of interest here, they are as yet unable to probe the small mixing-angle regimes which are important in supernovae (Cowsik & McClelland 1972; see also Harari 1989, and references therein). Furthermore it is not clear whether a ν_μ-ν_τ mass splitting in the range of $\Delta^{1/2}$ discussed here implies a similar ν_e-ν_τ splitting. Harari (1989) has outlined a proposal for a $\nu_\mu \rightleftharpoons \nu_\tau$ oscillation experiment which would probe closure mass $\Delta^{1/2}$ values and be sensitive down to vacuum mixing angles of $\sin^2 2\theta > 4 \times 10^{-4}$, which is at the upper end of the range which would be interesting in supernova shock reheating.

6. DISCUSSION AND CONCLUSIONS

We have discussed how a ν_μ or ν_τ neutrino with a cosmologically important mass (10–100 eV) would give a neutrino mass level crossing with a light ν_e in a region between the neutrino sphere and the shock in postbounce Type II supernova models. Significant neutrino transformation at this resonance may yield a harder spectrum of ν_e with a higher flux behind the shock which may, in turn, lead to enhanced matter heating rates. Our numerical calculations indicate that a 60% or so increase in the supernova energy would be expected. This may be of significance in delayed-mechanism supernova models which are at this point still fraught with uncertainty due to the difficulty of dynamic neutrino transport calculations.

We note that Woosley & Hoffman (1992) and Meyer et al. (1991) have discussed how the "hot bubble" which forms in the late stages of Type II supernova models may provide a near-perfect environment for the *r*-process, as the entropy is high, the overall density is low, and the free-neutron number density is appreciable. However, there is a question regarding whether or not the entropy in the supernova models at early times is high enough to effect a good *r*-process. The extra heating discussed in this paper from putative neutrino oscillations would give an added boost to the rate at which the entropy comes up as the hot bubble is formed. The entropy should be boosted by a factor of order the ratio of the extra energy deposited to the temperature: in other words, factors of 2–3. This may be quite important for *r*-process models. The added effect of neutrino-spallation (Woosley et al. 1990). will be to smooth the *r*-process abundance distribution.

Finally we have not discussed whether the ν_x which would mix with the ν_e in the reheating region is a ν_μ or ν_τ. Likely it would be the ν_τ. In this case the standard "seasaw" models of neutrino mass hierarchies could be stretched to accommodate a closure mass ν_τ and a ν_e-ν_μ vacuum mass splitting in a range adequate to give a level crossing in the Sun (cf. Bethe 1986). If the resolution of the solar neutrino problem turns out to have nothing to do with neutrino oscillations, then, of course, this argument is specious. Yet another consideration is the persistent report of evidence for a 17 keV neutrino (cf. Simpson 1985; Norman 1991) which mixes with the ν_e, as inferred from the beta spectrum distortions in several beta-unstable nuclides. If this evidence holds up, and the neutrino in question is identified with the ν_τ, then a level crossing between ν_e and ν_μ in the Sun would leave us with no appropriate level crossing in the supernova re-heating region. The lack of firm laboratory evidence for, or against, massive neutrinos with flavor mixings highlights the importance of astrophysical considerations.

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