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Abstract

We have calculated the thermal flux of low-energy solar neutrinos and antineutrinos of all flavors. The principal process, neutrino pair emission from free-bound atomic transitions in iron and other solar metals, produces a peak flux density per flavor $\sim 10^{11}/\mathrm{cm}^2/\mathrm{sec}/\mathrm{MeV}$ at $E_\nu \sim 1$ keV, a value comparable to the peak of the pp electron neutrino density. This flux carries information on the temperature and metallicity of the solar core and on heavy flavor neutrino masses for m_{ν_μ} or $m_{\nu_\tau} \gtrsim 1$ keV.

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A great deal of effort has been invested in detecting the solar electron neutrinos produced as a byproduct of solar fusion [1]. This letter focuses on another source of solar neutrinos, those produced by charged- and neutral-current thermal processes operating in the solar core. While the mean energy of these neutrinos is low, on the order of the solar core temperature $kT_c \sim 1.3$ keV, their flux density at earth is quite remarkable, $\sim 10^{11}/\mathrm{cm}^2/\mathrm{sec}/\mathrm{MeV}$ per flavor. Thus, above microwave energies $E_\nu \gtrsim 10^{-2}$ eV, this flux is the dominant source of antineutrinos of all flavors and of muon and tauon neutrinos. In this letter we describe the processes producing such neutrinos and comment on possibilities for their detection.

Various mechanisms for generating pair neutrinos are illustrated in Fig. 1. The first, Compton production, has been evaluated with the inclusion of the plasmon pole. The second is the decay of a plasmon into a $\nu\bar{\nu}$ pair. Both the Compton and plasmon processes are familiar neutrino cooling processes in stars [2]. The third is the decay of a thermally populated excited nuclear state by pair emission, a process that can contribute at solar temperature only in the case of an anomalously low energy M1 transition. Bahcall, Treiman, and Zee [3] identified the 14.4 keV Mossbauer transition in ^{57}Fe as the principal contributor in the sun. The fourth process is pair emission as an electron makes the transition from a free state to a bound atomic orbital. Free-bound transitions are known to be important to the solar opacity [4]. Although we know of no previous calculation of the analogous weak process, we find that free-bound transitions to atomic states in Fe and other metals completely dominate thermal neutrino production.

The evaluation of Compton pair production off free electrons is simplified by the fact that $T_c \ll m_e$, allowing one to treat the electrons in a static approximation. That is, they are allowed to absorb momentum but no energy. For the weak Hamiltonian

$$H = -\frac{G}{2\sqrt{2}}\bar{e}(g_V\gamma_\mu - g_A\gamma_\mu\gamma_5)e\ \bar{\nu}\ \gamma^\mu(1-\gamma_5)\nu\ , \tag{1}$$

the number of pairs produced per unit time per unit volume is found to be

$$\frac{dN}{dt} = \frac{1}{6\pi^4}\frac{\alpha G^2}{m_e^2}(g_V^2 + 5g_A^2)\left[\frac{2}{(2\pi)^3}\int d\vec{p}_i \frac{1}{(e^{(\epsilon_{p_i}-\mu)/kT}+1)}\left[1 - \frac{1}{(e^{(\epsilon_{p_f}-\mu)/kT}+1)}\right]\right]$$
$$\int_0^\infty \frac{1}{\omega_\gamma}\frac{1}{e^{\omega_\gamma/kT}-1}\epsilon_\nu^2\epsilon_{\bar{\nu}}^2(\omega_\gamma^2 - \tfrac{2}{3}\epsilon_\nu\epsilon_{\bar{\nu}})d\epsilon_\nu d\epsilon_{\bar{\nu}} \ , \tag{2}$$

where ϵ_ν and $\epsilon_{\bar{\nu}}$ are the neutrino and antineutrino energies, ω_γ the photon energy, and ϵ_{p_i} and ϵ_{p_f} the initial and final electron energies. The static electron approximation identifies $\omega_\gamma = \epsilon_\nu + \epsilon_{\bar{\nu}}$ and $\epsilon_{p_i} = \epsilon_{p_f} = \vec{p}_i^{\,2}/2m_e$. The chemical potential μ is a function of position within the sun, and has been calculated from the temperature and density distributions of the standard solar model, as given by Bahcall [5]. In the solar core, $\mu \sim -2.1$ keV, so that the electron gas is only slightly degenerate. For heavy flavor neutrinos $g_V = 1 - 4\sin^2\theta_w$ and $g_A = 1$, while for electron neutrinos one obtains, after a Fierz transformation of the charged current interaction, the effective couplings $g_V = -(1 + 4\sin^2\theta_w)$ and $g_A = -1$. Thus the ratio of the heavy flavor and electron neutrino rates is ~ 0.58. The contribution from Compton scattering off ionized nuclei is down by a factor of order $(m_e/M_N)^2$.

This result is modified by the plasmon pole contributions of Fig. 1b, which we believe have not been evaluated previously. Transverse plasmons satisfy the dispersion relation

$$\omega_\gamma^2 = \vec{q}^{\,2} + \omega_{p\ell}^2\left(1 + \frac{\vec{q}^{\,2}}{\omega_\gamma^2}\frac{T}{m_e}\right) \ , \tag{3}$$

where the plasma frequency $\omega_{p\ell} = \sqrt{\frac{4\pi\alpha n_e}{m_e}} \sim 0.28$ keV for an electron number density n_e characteristic of the sun's center. Longitudinal plasmons (charge density waves) satisfy a dispersion relation

$$\omega_\gamma^2 = \omega_{p\ell}^2\left(1 + \frac{3\vec{q}^{\,2}}{\omega_\gamma^2}\frac{T}{m_e}\right) \ . \tag{14}$$

Note that $\frac{T}{m_e} \sim v_e^2 \sim 0.005$. Thus transverse plasmons propagate as photons of mass $\omega_{p\ell}$, while longitudinal plasmons carry an energy equal to the plasma frequency. For

neutrinos of present interest ($\epsilon_\nu \sim$ keV) the transverse plasmon pole will be encountered in the angular integrations: the q_μ^2 carried off by the neutrino pair is timelike and can be made small ($\sim \omega_{p\ell}^2$) by letting the opening angle between the neutrinos approach zero. In contrast, the longitudinal plasmons become important only when $\epsilon_\nu + \epsilon_{\bar{\nu}} \sim \omega_{p\ell}$.

The evaluation of the pole diagrams is somewhat complicated and, for that reason, will be presented elsewhere [6]. Diagrams involving axial couplings of the weak bosons to electrons are found to be negligible, suppressed by $(\alpha \frac{T}{m_e})^2$ in the rate. While the simple Compton and pole contributions interfere, this interference can be ignored since the pole contribution is so narrowly peaked at kinematics defined by $q_\mu^2 \sim \omega_{p\ell}^2$. It follows that pole contribution is proportional to g_V^2 and thus is significant only for electron neutrinos. The plasmon widths have been taken from the two-particle two-hole calculation of DuBois *et al.* [7]. (Decay into one-particle one-hole electron states is forbidden for timelike q_μ^2.) The terrestrial neutrino flux densities produced by the pole and simple Compton processes are given in Fig. 2. They were determined by folding Eq. (2) and the corresponding pole result with the solar density and temperature profiles of Ref. [5].

The contribution of direct transverse plasmon decays into $\nu\bar{\nu}$ pairs (Fig. 1c) is also shown in Fig. 2. The neutrino pair production rate per unit volume is

$$\frac{dN}{dt} = \frac{G^2}{128\pi^4} \frac{g_V^2}{\alpha} \omega_{p\ell}^6 \int \frac{1}{\omega_\gamma^2} \frac{1}{e^{\omega_\gamma/kT} - 1} (\epsilon_\nu^2 + \epsilon_{\bar{\nu}}^2) d\epsilon_\nu d\epsilon_{\bar{\nu}} \tag{5}$$

where $\omega_\gamma = \epsilon_\nu + \epsilon_{\bar{\nu}}$ and the integrations extend over all neutrino energies satisfying $\omega_\gamma > \omega_{p\ell}$. We have ignored the contribution of longitudinal plasmons, which again contribute only if $\epsilon_\nu + \epsilon_{\bar{\nu}} \sim \omega_{p\ell}$.

Neutral current decays of excited nuclear states can produce pair neutrinos (Fig. 1d). The importance of a given element depends on the probability that the excited state will be thermally populated, on the element's solar mass fraction, and on the strength

of the isovector axial (M1) transition matrix element. As noted in Ref. [3], the important metal is ^{57}Fe because of the low-lying $3/2^- \rightarrow 1/2^-$ Mössbauer transition. We can express the pair neutrino rate per flavor in terms of the known gamma decay rate $\omega_\gamma = 7.41 \cdot 10^5$/sec (thereby eliminating most of the nuclear physics uncertainty)

$$\frac{dN}{dt} = \frac{\omega_\gamma}{1+\delta^2} N(^{57}Fe) \frac{3}{\pi^3} \frac{G^2 F_A^2}{\alpha} \frac{M_N^2}{W_0^3} \frac{e^{-W_0/T}}{[\mu_1 + (\mu_0 - \frac{1}{2})\beta - \eta]^2} \int_0^{W_0} \epsilon_\nu^2 (W_0 - \epsilon_\nu)^2 \, d\epsilon_\nu \, , \tag{6}$$

where $N(^{57}\mathrm{Fe})$ is the number density of ^{57}Fe nuclei in the solar core, $W_0 = 14.4$ keV is the $3/2^- \rightarrow 1/2^-$ transition energy, $\delta = 0.0022$ is the E2/M1 mixing ratio, $\mu_0 = 0.88$ and $\mu_1 = 4.71$ are the isoscalar and isovector magnetic moments, and $F_A = 1.26$ is the axial vector coupling constant. The nuclear-structure-dependent terms

$$\eta = -\frac{< J_f || \sum_{i=1}^A \vec{\ell}(i)\tau_3(i) || J_i >}{< J_f || \sum_{i=1}^A \vec{\sigma}(i)\tau_3(i) || J_i >} \sim 0.80 \tag{7a}$$

and

$$\beta = \frac{< J_f || \sum_{i=1}^A \vec{\sigma}(i) || J_i >}{< J_f || \sum_{i=1}^A \vec{\sigma}(i)\tau_3(i) || J_i >} \sim -1.19 \tag{7b}$$

are taken from the shell model calculations of Ref. [8]. Note that the maximum rate for Eq. (6) occurs when $W_0 = 5T$, or about 6.5 keV in the solar core, so the ^{57}Fe transition would be more effective in somewhat hotter stars. The neutrino flux densities in Fig. 2 are obtained by integrating over the solar density/temperature profile, assuming the ^{57}Fe solar mass fraction $X = 3.26 \cdot 10^{-5}$ [9].

The final neutrino pair process we consider is Z^0 emission in the transition of a continuum electron to a bound atomic orbital (Fig. 1e). Although such free-bound transitions to atomic states in Fe and other metals are known to be an important contributor to the solar opacity [4], we are unaware of any calculation of the analogous weak process.

The rate for capturing into Bohr orbits described by (n, ℓ) (summed over all magnetic quantum numbers) is

$$\frac{dN}{dt} = N(Z)\, f(n,\ell)\, \frac{8\sqrt{2}}{\pi^4}\, m_e^{3/2} G^2 (g_V^2 + 3g_A^2) \left(\frac{na_0}{Z}\right)^3 (2\ell+1) 2^{4\ell} (\ell!)^2$$
$$\frac{n(n-\ell-1)!}{(n+\ell)!} \int \epsilon_\nu^2 d\epsilon_\nu \epsilon_{\bar\nu}^2 d\epsilon_{\bar\nu} \frac{\sqrt{\epsilon_p}}{(e^{(\epsilon_p-\mu)/kT}+1)} \frac{\xi^{2\ell}}{(1+\xi^2)^{2\ell+4}} \left[C_{n-\ell-1}^{\ell+1}\left(\frac{\xi^2-1}{\xi^2+1}\right)\right]^2 , \tag{8}$$

where $\epsilon_p = \epsilon_\nu + \epsilon_{\bar\nu} - E_b(n,\ell)$ is the continuum electron's energy, $E_b(n,\ell)$ is the (positive) Bohr orbit binding energy, $a_0 = \hbar^2/m_e e^2$ is the hydrogen Bohr radius, Z is the atomic charge, $N(Z)$ is the number density of the atoms in the solar core, the $C_{n-\ell-1}^{\ell+1}$ are Gegenbauer polynomials arising from the Fourier transform of the atomic wave functions [10], and $f(n,\ell)$ is a blocking factor correcting for the incomplete ionization of the atomic shells. (In the solar interior, the only atomic shells not completely ionized are those with low n and high Z.) The integrations over ϵ_ν and $\epsilon_{\bar\nu}$ are restricted by the requirement $\epsilon_p > 0$. In deriving Eq. (8) the integration over neutrino angles was performed by making the approximation

$$\xi = \frac{na_0}{Z} |\vec{p}_\epsilon - \vec{p}_\nu - \vec{p}_{\bar\nu}| \sim \frac{na_0 p_e}{Z}$$

as $p_e \sim \sqrt{2m_e(\frac{3}{2}kT)} \sim 45$ keV $\gg \epsilon_\nu$.

The qualitative features of free-bound transitions can be understood readily. As $\xi \sim 12n/Z \gg 1$ for light metals, we see that higher ℓ contributions are suppressed by a factor $1/\xi^{2\ell}$. The dominant $\ell = 0$ contributions vary as $1/n^3$, so that one would expect the $n = 1$ state to carry $\approx 80\%$ of the strength. The $n = 1$ capture rate varies, for small Z, as Z^5: the more deeply bound wave functions of higher Z atoms begin to contain the high momentum components that can match p_e.

Neglecting the blocking factor, one can maximize the $n = 1$ capture rate of Eq. (8) as a function of Z. One finds $Z_{max} = \sqrt{\frac{5}{3}} p_e a_0 \sim 16$, using $p_e \sim 45$ keV. (A somewhat

higher p_e would be more appropriate, as electrons on the high energy tail of the thermal distribution are favored by the phase space factors in Eq. (8).) In our numerical calculations we find, using the H, He, and CNO distributions from Ref. [5], Cameron abundances for heavier metals, and summing over all (n, ℓ) until results are fully converged, that Fe is the dominant metal, followed by O, S, and Si. The resulting terrestrial $\nu_e/\bar{\nu}_e$ flux densities, found by integrating Eq. (8) over the solar composition and temperature profiles [5], are shown in Fig. 3. The corresponding heavy-flavor flux densities are smaller by ~ 0.46. Free-free and bound-bound transitions also occur, but the rates are small compared to free-bound ones.

In Fig. 4 we summarize all of the neutrino sources discussed in this letter. The peak of the thermal solar ν_e (or $\bar{\nu}_e$) flux density, occurring at ~ 1 keV, is $\sim 10^{11}/\mathrm{cm}^2/\mathrm{sec}/\mathrm{MeV}$, a value nearly identical to the peak of the pp fusion neutrino flux density. The heavy flavor fluxes (ν_μ, $\bar{\nu}_\mu$, ν_τ, or $\bar{\nu}_\tau$) are about a factor of two smaller. The integrated ν_e and ν_μ fluxes are $2.4 \cdot 10^8/\mathrm{cm}^2/\mathrm{sec}$ and $1.1 \cdot 10^8/\mathrm{cm}^2/\mathrm{sec}$, respectively.

The solar thermal neutrinos are a remarkable feature of the astrophysical neutrino flux passing through the earth. The thermal $\bar{\nu}_e$ and heavy flavor fluxes are the dominant source of these flavors above cosmic background neutrino energies, $E_\nu \gtrsim 10^{-2}$ eV (see Fig. 4). The thermal ν_es are the dominant flux density from 10^{-2} eV to 10 keV, where the pp neutrinos take over.

The thermal neutrinos are unimportant to solar evolution as they account for only 0.1% of solar neutrino energy loss. However this flux carries a great deal of information about the solar model and about neutrino properties. The production mechanisms depend on well-understood atomic processes, and the resulting flux depends linearly on the core metallicity Z. Thus a measurement of this flux could test an important assumption of the standard solar model [5], that core abundances for $Z > 5$ at the onset of the main sequence can be equated to today's surface abundances. The high

energy tail of the thermal flux is quite sensitive to the core temperature: a reduction in T_c by 5% (often postulated in nonstandard solar models producing reduced ^{8}B fluxes) diminishes fluxes by $\sim 43\%$ at $\epsilon_\nu \sim 10$ keV. An appreciable ν_τ flux will be produced only if $m_{\nu_\tau} \lesssim$ several kT: the current mass limit is 31 MeV. Finally, these neutrinos probe the small-δm^2 region of the MSW δm^2-$\sin^2 2\theta$ plane [11], including $\bar{\nu}_e \leftrightarrow \bar{\nu}_\mu$ oscillations that would arise for an inverted mass hierarchy.

Is this flux simply a curiosity, or is there some prospect for detecting thermal neutrinos? The most intriguing possibility is the resonant capture reaction

$$\bar{\nu}_e + e^- + (A, Z) \rightarrow (A, Z-1),$$

a process that depends on the $\bar{\nu}_e$ flux density at the resonant energy, and that has been previously discussed in connection with terrestrial $\bar{\nu}_e$ sources [12]. Cross sections could thus be very large in high Z atoms, where the electron overlap with the nucleus is favorable. There appear to be a few cases involving nuclear transition energies of 2-20 keV. We will present the obstacles to such a measurement in a future paper [6].

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Figure Captions

Figure 1: Representative diagrams of thermal processes for neutrino pair production: a) Compton process; b) plasmon pole contribution to the Compton process; c) transverse plasmon decay; d) nuclear Z^0 emission; and e) pair production in free-bound atomic transitions.

Figure 2: Neutrino flux densities from free-bound transitions (long dashed line), the Compton pair process (short dashes), the pole contribution to the Compton pair process (dashed-double-dot), ^{57}Fe nuclear decay (dashed-dot), and transverse plasmon decay (dotted). Totals (solid line) are for a single flavor.

Figure 3: The total free-bound ν_e flux density (solid line) and the separate contributions of the most important metals (Fe, S, Si, O; dashed lines) are shown as a function of the neutrino energy. The contributions from H and He are also shown. The $\bar{\nu}_e$ results are identical, and the heavy-flavor neutrino fluxes are about a factor of two smaller.

Figure 4: Natural neutrino sources. The terrestrial $\bar{\nu}_e$ flux and continuous flux of extragalactic supernova neutrinos of all flavors are from Ref. [12]. The solar (fusion) ν_e flux is the standard solar result of Ref. [3]. The thermal solar neutrinos are for a single flavor.

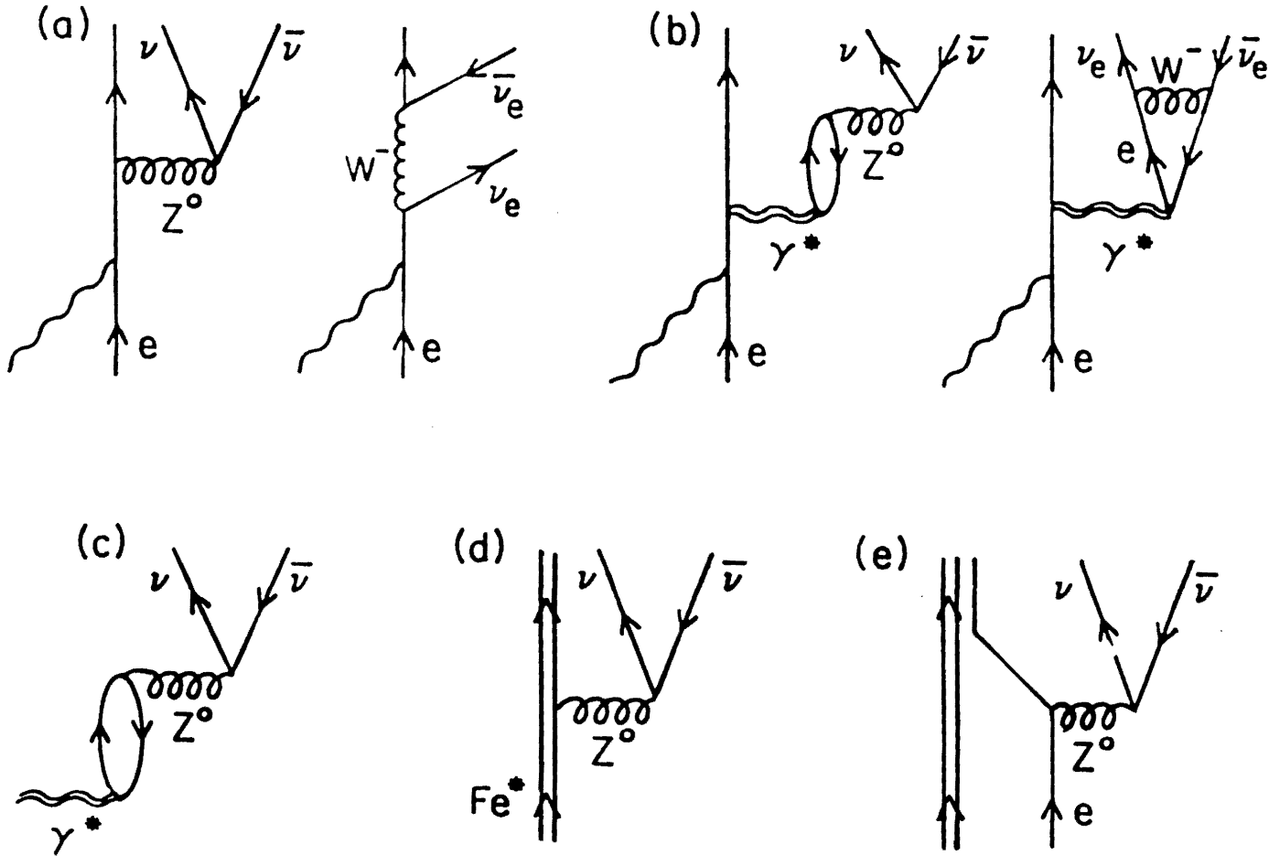

Figure 1

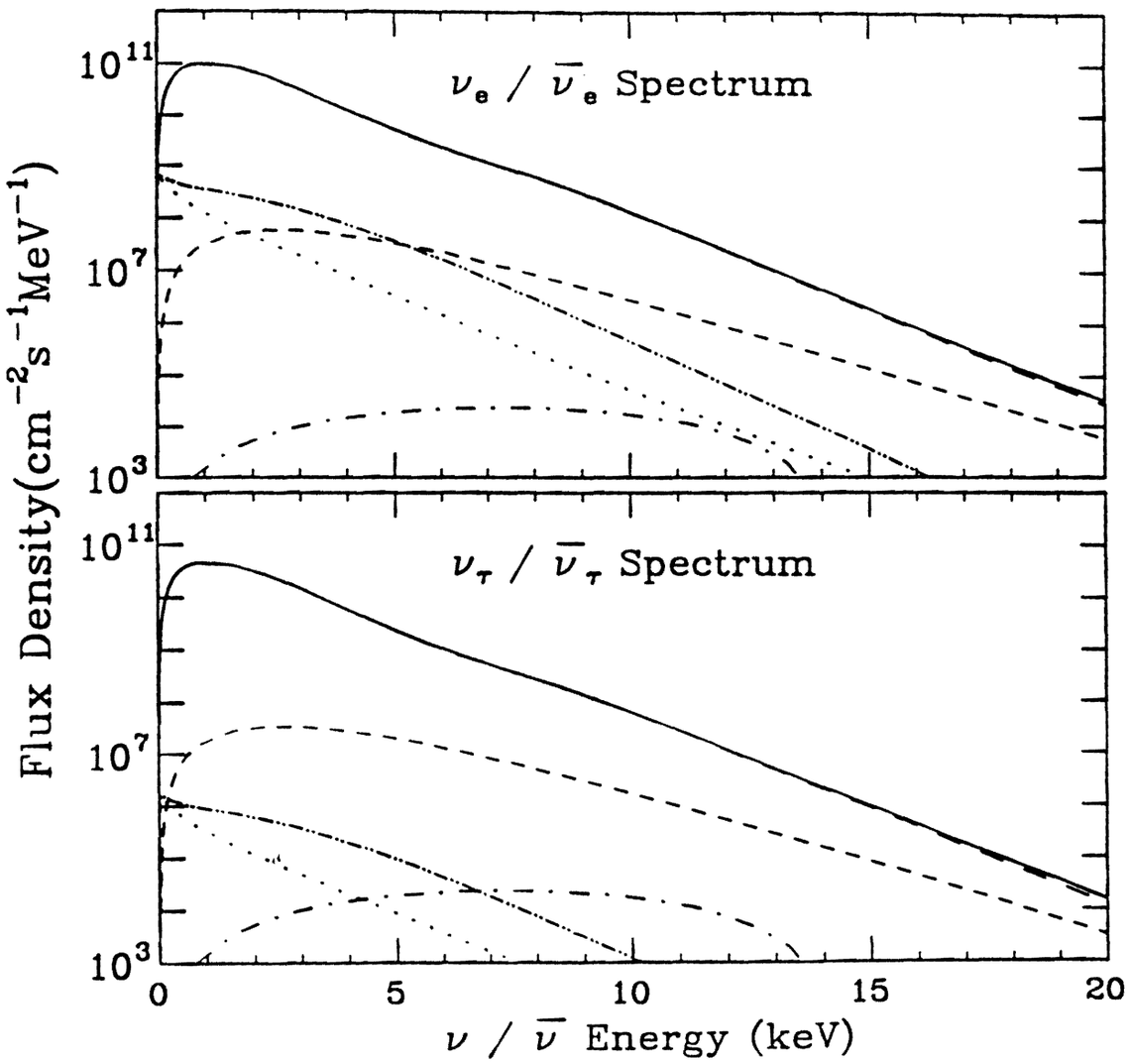
ν_e / $\bar{\nu}_e$ Spectrum
ν_τ / $\bar{\nu}_\tau$ Spectrum
Flux Density($cm^{-2}s^{-1}MeV^{-1}$)
10^{11}
10^{7}
10^{3}
0
5
10
15
20
ν / $\bar{\nu}$ Energy (keV)

Figure 2

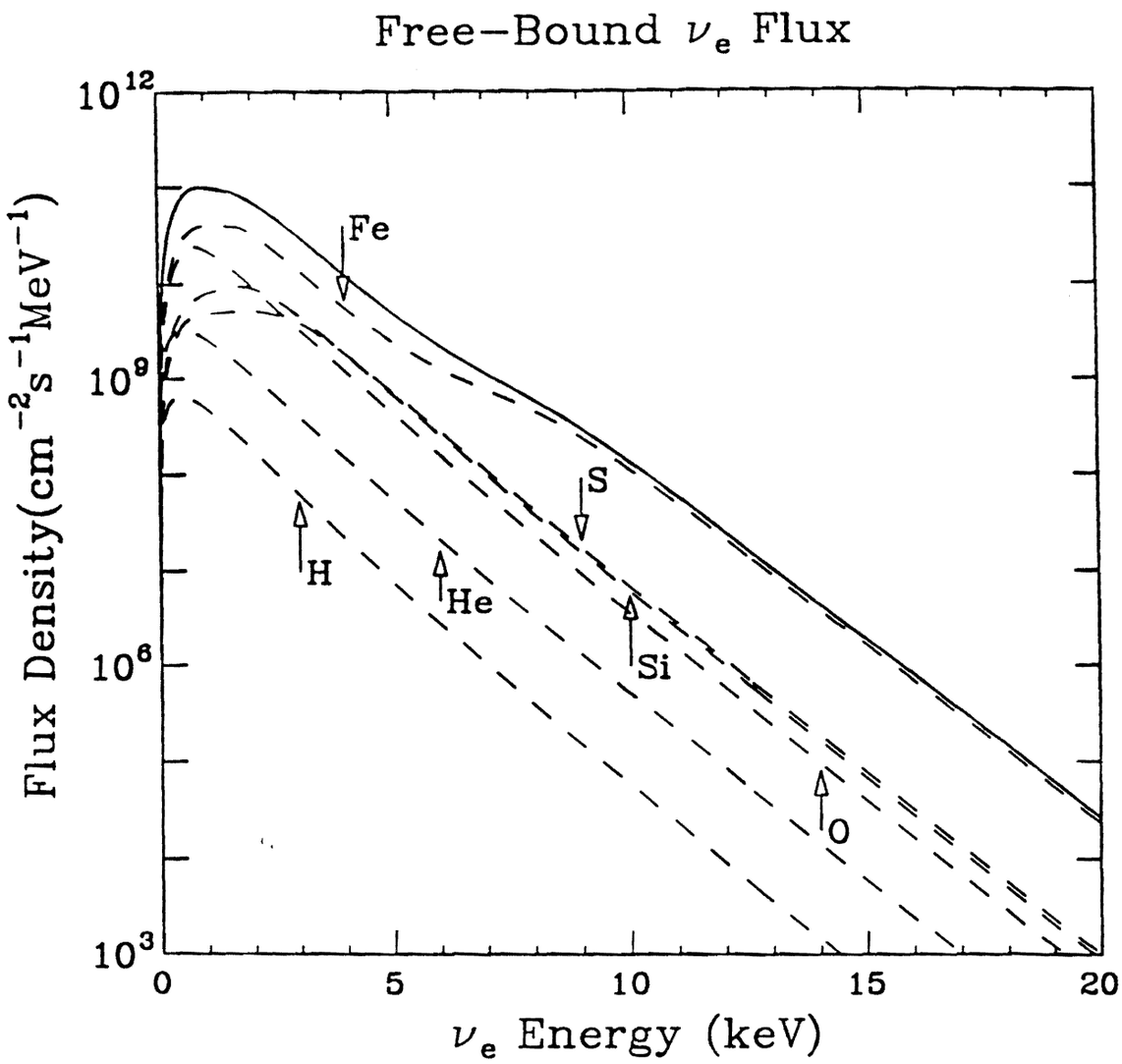
Free–Bound ν_e Flux
10^{12}
10^{9}
10^{6}
10^{3}
Flux Density($cm^{-2}s^{-1}MeV^{-1}$)
Fe
S
H
He
Si
O
0
5
10
15
20
ν_e Energy (keV)

Figure 3

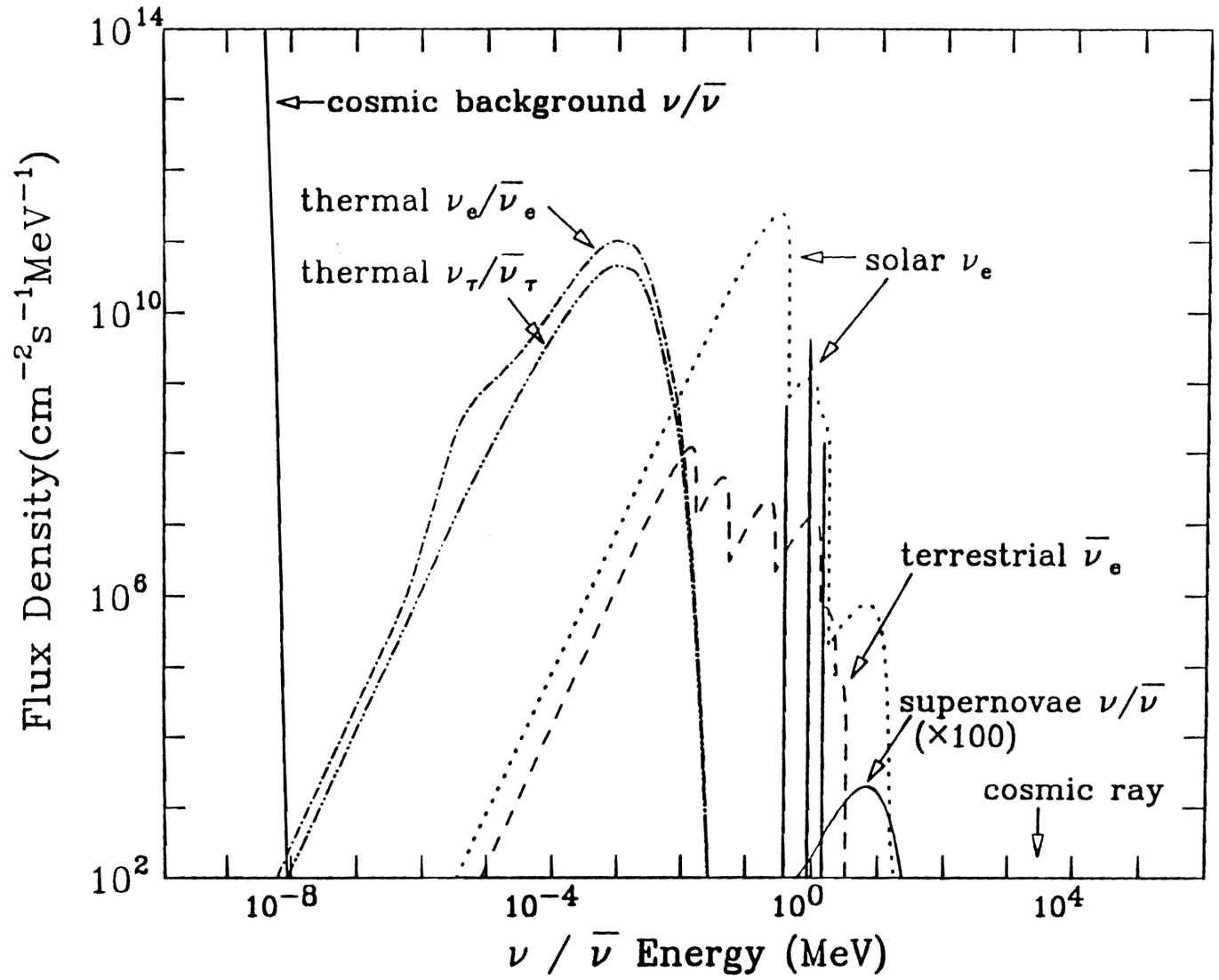
cosmic background $\nu/\bar{\nu}$
thermal $\nu_e/\bar{\nu}_e$
thermal $\nu_\tau/\bar{\nu}_\tau$
solar ν_e
terrestrial $\bar{\nu}_e$
supernovae $\nu/\bar{\nu}$ (×100)
cosmic ray
Flux Density(cm^{-2}s^{-1}MeV^{-1})
10^{14}
10^{10}
10^{6}
10^{2}
10^{-8}
10^{-4}
10^{0}
10^{4}
ν / $\bar{\nu}$ Energy (MeV)

Figure 4