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Decay of reactor neutrinos

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We consider the decay of massive neutrinos which couple to electrons and are, therefore, produced in nuclear reactors. Lifetime limits for the γ and electron-positron decay modes of these neutrinos are deduced from the experimental limit on the singles count rate in the detector used to study neutrino oscillations at the Gösgen reactor. The dominantly coupled neutrinos are light, and their invariant-lifetime limit $t^{\text{c.m.}}/m_\nu$ is 1—3 sec/eV. The subdominantly coupled heavy neutrinos with mass 1—4 MeV could decay into electron-positron pairs. These pairs were not observed, and from the absence of such a signal we deduce restrictions on the corresponding mixing parameters.

If neutrinos have mass, the heavier ones could decay into the lighter ones. Indeed, neutrino decay is an inevitable consequence of the neutrino mass and mixing, and the search for the decay is therefore an important part of the study of these phenomena. In this note the limits on the rate of the decay modes

$$\nu_2 \to \nu_1 + \gamma \tag{1}$$

and

$$\nu_2 \to \nu_1 + e^+ + e^- \tag{2}$$

obtained in experiments on nuclear reactors are derived.

The neutrino decay modes (1) and (2) were considered in several papers,[1–4] where the decay rates are evaluated in terms of the neutrino masses and mixing angles. (The existing data on the neutrino decay are summarized in Ref. 5.)

In order to connect the decay rate of the initial neutrino at rest with the count rate in a detector, it is necessary to consider the kinematics of the decay. For the reactor neutrinos this is done here for moving initial neutrinos and a detector which counts only photons or electron-positron pairs with certain selected energies. Using these results we discuss the lifetime limits of the decays (1) and (2) which are derived from limits on the singles count obtained in the course of the neutrino-oscillation searches.[6–8]

Throughout, we consider a beam of neutrinos ν_2 with mass m_{ν_2}. The mass of the final fermion ν_1 is assumed to be negligibly small; ν_1 could, in fact, be massless. Nuclear reactors produce neutrinos ν_2 if these ν_2's couple to electrons and if m_{ν_2} is less than the characteristic energy ($\sim$ several MeV) of the nuclear β decay.

Two situations should be distinguished. If ν_2 is dominantly coupled to electrons, it is known that $m_{\nu_2} \leq 60$ eV (or, perhaps, $m_{\nu_2} \simeq 30$ eV). Reactor-produced ν_2 are then ultrarelativistic and their spectrum is unaffected by their mass; it is just the familiar reactor antineutrino spectrum.[9] Only the photon decay mode (1) could be present. On the other hand, ν_2's could be also relatively heavy and subdominantly coupled to electrons. The spectrum then depends on m_{ν_2} and on the mixing parameter U_{e2}. If $m_{\nu_2} > 2m_e$ the electron-positron decay (2) becomes the main decay channel.

Let us consider first the γ decay (1) of the dominantly coupled light initial antineutrinos ($U_{e2} \simeq 1$). For a reactor with R fissions per second and a detector of volume V at a distance d from the reactor core there are

$$dN(E_\nu) = \frac{RV}{4\pi d^2} \frac{n(E_\nu)}{c} dE_\nu \tag{3}$$

ν_2 neutrinos with the energy E_ν in the detector. Here $n(E_\nu)$ is the continuous antineutrino spectrum associated with fission[9] and c is the velocity of light.

The energy of the decay photon depends on the initial neutrino energy and on the center-of-mass angle θ between the photon momentum and the beam direction,

$$E_\gamma = \frac{E_\nu}{2} \left[1 + \frac{p_\nu}{E_\nu} \cos\theta \right] \simeq \frac{E_\nu}{2} (1 + \cos\theta) \ . \tag{4}$$

Consequently, the photon laboratory energy spectrum depends on the c.m. angular distribution, which is generally of the form

$$dN = \tfrac{1}{2}(1 + a \cos\theta) d \cos\theta \ , \tag{5}$$

where $|a| \leq 1$. The angular anisotropy is the result of parity nonconservation in the decay and of the nonvanishing polarization of the neutrinos. For Majorana neutrinos one has, clearly, $a=0$, while for Dirac neutrinos and left-handed coupling[10] $a=-1$.

Combining Eqs. (3)—(5) and taking into account the relativistic time dilatation one obtains the laboratory photon spectrum

$$\frac{dN}{dE_\gamma} = m_\nu \Gamma^{\text{c.m.}} \int_{E_\gamma}^{\infty} \frac{1 + a - 2aE_\gamma/E_\nu}{{E_\nu}^2} N(E_\nu) dE_\nu \ , \tag{6}$$

where $\Gamma^{\text{c.m.}}$ is the center-of-mass decay rate. The integral in Eq. (6), with $N(E_\nu)$ replaced by $n(E_\nu)$ to make it independent of the specific reactor and detector, is shown in Fig. 1 as a function of E_γ. From Eq. (6) and the experimental photon count rate one obtains a value (or a limit) of $m_\nu \Gamma^{\text{c.m.}}$ (or equivalently $t^{\text{c.m.}}/m_\nu$).

In the only existing experiment dealing with electron neutrinos the laboratory lifetime limit of 6×10^7 sec was established.[11] However, this result was obtained assuming that all antineutrinos should be counted while only photons with energies 0.1–0.5 MeV were detected. When the analysis is performed in the way described here, with the antineutrino flux, detector volume, and photon count rate of Ref. 11, one arrives at the invariant-lifetime limit $t^{\text{c.m.}}/m_\nu > 30$ sec/eV, poorer than the result quoted by the Particle Data Group.[12] (Note that the main motivation of Ref. 11, the proof that the neutrino decay length far exceeds the sun-Earth distance, remains certainly true.) For muon neutrinos the best existing limit[13] is $t^{\text{c.m.}}/m_\nu > 9$ sec/eV.

Here we derive neutrino-lifetime limits from the photon singles count rate obtained during the neutrino oscillation search at Gösgen.[7,8] The singles count rate reactor on minus reactor off is compatible with zero and restricted to 2 counts/sec (1σ) for the discriminator threshold of 0.6 MeV in each of the target cells of the detector. (For 2-MeV threshold one obtains the limit of 0.5 counts/sec.)

The detector at $d = 38$ m is characterized by the quantity

$$RV/4\pi d^2 c = 10^7 .$$

The photon detection efficiency has been estimated by a crude Monte Carlo calculation. The efficiency is essentially constant, $\epsilon \simeq 0.33$, for $1 \le E_\gamma \le 4$ MeV, and it decreases to zero when E_γ approaches the threshold value.

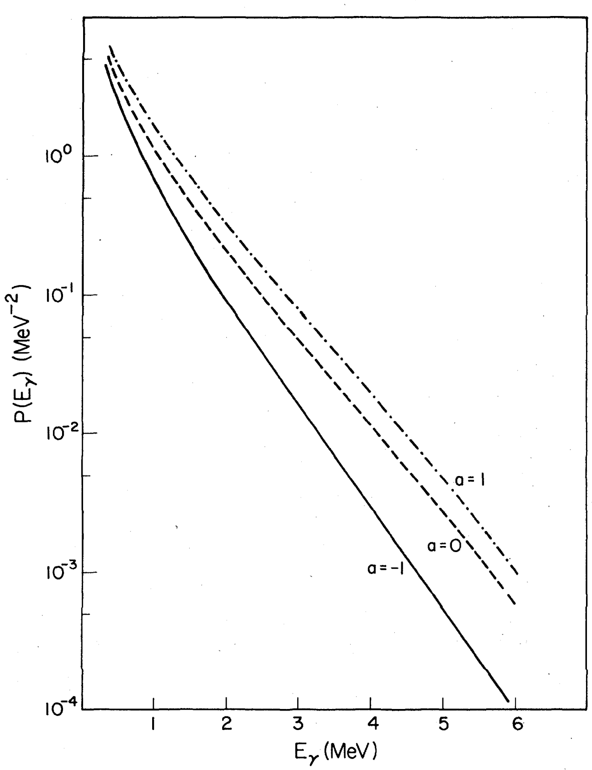

FIG. 1. The photon spectrum $P(E_\gamma)$, Eq. (6), for different values of the angular-distribution parameter a.

From the graphs in Fig. 1, the estimated efficiency, and from the experimental count rates for the 0.6-MeV threshold we obtain the following lifetime limits in sec/eV

$$\frac{t^{\text{c.m.}}}{m_\nu} \ge 1 \ (a=-1),\ 1.8\ (a=0),\ 2.5\ (a=1) . \tag{7}$$

Consistent (although somewhat less restrictive) limits are obtained using the 2-MeV threshold at Gösgen or the singles counts at Institut Laue-Langevin.[6]

All of the quoted laboratory lifetime limits are of a similar magnitude, nowhere near the theoretical estimates based on the "standard" assumptions[1–5] (Dirac neutrinos, three known lepton generations). However, it is useful to establish rigorous experimental lifetime limits since one cannot exclude the existence of very heavy charged leptons or other mechanisms which would speed up the decay.

Next we consider the electron-positron decay (2) of the subdominantly coupled heavy ($m_{\nu_2} > 2m_e$) neutrinos. Equation (3) must be now multiplied by $|U_{e2}|^2$ to obtain the number of antineutrinos $\bar{\nu}_2$ in the detector. Note that the ν_2's may be nonrelativistic, $v_\nu < c$; however, the factor c/v_ν in the dwelling time of the neutrinos in the detector is canceled by the factor v_ν/c in the phase space of the β decay involving heavy neutrinos.

In the reactor experiments one measures the total laboratory energy E of the electron-positron pair. (The kinetic energy is measured with essentially 100% efficiency, the annihilation radiation has rather small detection efficiency in a given target cell.) For a given energy E_ν and mass m_ν of the initial neutrino ν_2 the energy E is restricted to the interval

$$E_\nu \left[1 - \frac{1+v_\nu}{2} \left(1 - \frac{4m_e^2}{m_\nu^2} \right) \right] \le E \le E_\nu . \tag{8}$$

We have to consider again the c.m. angular distribution. Instead of the $e^- + e^+$ pair we may use the final neutrino ν_1 with the energy ϵ and emission angle θ. The c.m. double-differential distribution can be deduced from the formulas of Shrock[14] for Dirac neutrinos and the standard $V-A$ couplings,

$$\frac{d^2\Gamma_0}{d\epsilon\, d\cos\theta} = \Gamma_0(m_\nu)\,|U_{e2}U_{e1}|^2 (f_1 - \cos\theta f_s) , \tag{9}$$

where f_1 and f_s depend on the momentum transfer q^2 and on m_e/m_ν. The Γ_0 is independent of the energy and angle. The mixing angle U_{e2} in Eq. (9) is the same one as in the source term discussed above. On the other hand, it is natural to assume that the light neutrino ν_1 will be dominantly coupled to electrons and therefore $U_{e1} \sim 1$.

The decay rate in terms of E is of the form

$$\Gamma = \frac{\Gamma_0(m_\nu)\,|U_{e2}|^2}{E_\nu} \int_{E_{\min}}^{E_\nu} dE \int_{(\cos\theta)_{\min}}^{1} d\cos\theta \frac{f_1 - \cos\theta f_s}{1 + v_\nu \cos\theta} , \tag{10}$$

where $E_{\min}$ is the lower bound of E in Eq. (8) and

$$1 + v_\nu \cos\theta_{\min} = \frac{2(E_\nu - E)}{E_\nu(1 - 4m_e^2/m_\nu^2)} . \tag{11}$$

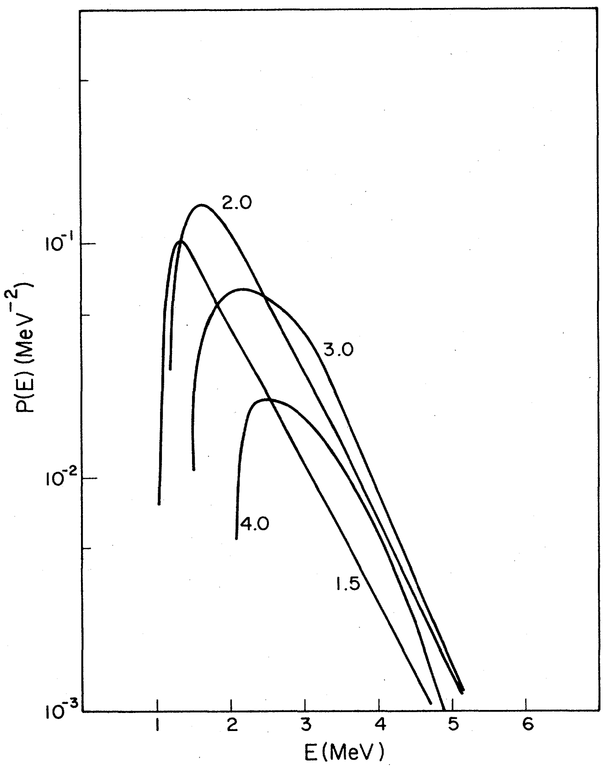

FIG. 2. The electron-positron spectrum [the integral in Eq. (12)]. The curves are labeled by the m_ν value in MeV.

Denoting the integral over $\cos\theta$ in (10) by $Y(E,E_\nu,m_\nu)$ we obtain from Eqs. (3), (8), and (10) the expected laboratory spectrum of the electron-positron pairs

$$\frac{dN}{dE}=\frac{RV}{4\pi d^2 c}\,|\,U_{e2}\,|^4 \times m_\nu \Gamma_0(m_\nu)\int_{E_{\nu\,\min}}^{E_{\nu\,\max}}\frac{n(E_\nu)}{E_\nu^{\,2}}Y(E,E_\nu,m_\nu)dE\ , \quad (12)$$

where the integration limits are

$$E_{\nu\,\min}=\max(m_\nu,E)\ ,$$

and $E_{\nu\,\max}$ is the solution of the equation

$$E=E_{\nu\,\max}\left[1-\frac{1+v_\nu}{2}\left[1-\frac{4m_e^{\,2}}{m_\nu^{\,2}}\right]\right].$$

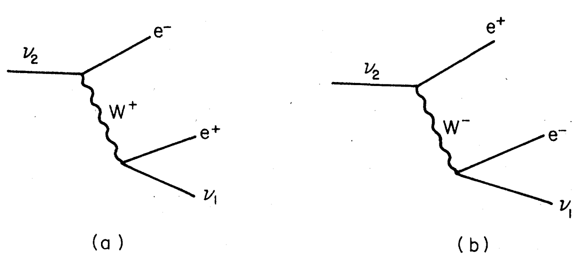

FIG. 3. Feynman graphs describing $\nu_2 \to \nu_1 + e^+ + e^-$. (a) is for Dirac neutrinos, (a) and (b) for Majorana neutrinos.

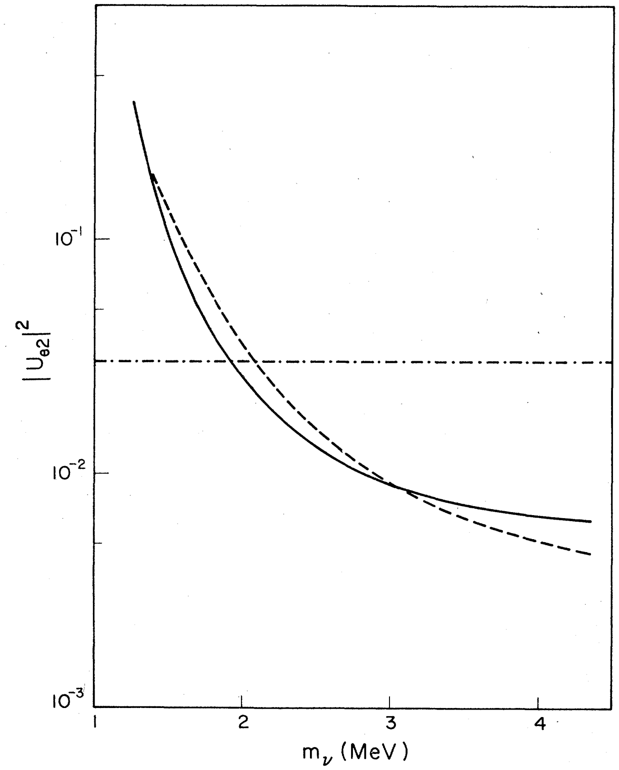

FIG. 4. Limits of the mixing parameter $|\,U_{e2}\,|^2$. Solid curve for threshold 0.6 MeV (2-counts/sec limit), dashed curve for threshold 2.0 MeV (0.5-counts/sec limit). The horizontal line is the 68%-C.L. limit from the neutrino-oscillation search (Ref. 7).

The integrals in Eq. (12) are shown in Fig. 2 for several m_ν values. The fast decrease at larger values of E is caused by the fast decrease of the reactor neutrino spectrum $n(E_\nu)$.

Theoretically, the c.m. decay rate for the mode (2) is obtained from the Feynman graph 3(a) for the Dirac neutrinos, and from Figs. 3(a) and 3(b) for the Majorana neutrinos. In the considered Dirac-neutrino case

$$\Gamma_0(m_\nu)=\frac{G_F^{\,2}}{192\pi^3}m_\nu^{\,5}\ , \quad (13)$$

where $G_F^{\,2}=0.206\,\mathrm{MeV}^{-5}\mathrm{sec}^{-1}$. Note that for the same mass of the initial neutrino the electron-positron decay mode (2) is expected to be considerably faster than the gamma decay mode (1). In fact, the experimental limit on the decay (2) allows one to deduce a meaningful restriction on the value of the mixing parameter U_{e2}. From the measured number of the electron-positron pairs (or from its limit), one thus obtains a value (or a limit) of $|\,U_{e2}\,|^4$ for each value of m_ν.

The singles count rate at Gösgen[6,7] quoted earlier represents also a limit on the kinetic energy of the positron electron pairs. To obtain a limit on $|\,U_{e2}\,|^4$ (respectively, $|\,U_{e2}\,|^2$), we integrate the differential spectrum in Eq. (12) from the threshold up. The resulting limits of $|\,U_{e2}\,|^2$ are shown in Fig. 4. The constant asymptotic value of $|U_{e2}|^2$ obtained in the neutrino-oscillation searches[6,7] is also shown. Thus, considerable improvement in the neutrino mass range $2<m_\nu<4$ MeV results. For $m_\nu>4$ MeV a better limit for the same mixing parameter is obtained from the study of the $\pi\to e+\nu$ decay.[15] Limits of the mixing parameter $|U_{\mu 2}|^2$ have been obtained recently by Minehart *et al.*[16] These limits, dealing with a different parameter of the neutrino mixing matrix, are for the $2<m_\nu<4$ MeV neutrino mass range approximately three times more restrictive than those in Fig. 4 for $|U_{e2}|^2$.

Very little is known about the mixing parameter $|U_{e2}|^2$ for the mass range 1—4 MeV. The only limits there come from the work of Toussaint and Wilczek[17] based on the expected production of heavy neutrinos in the sun, their decay into electron-positron pairs, and the experimental upper limit for the positron flux in the interplanetary space. The discrepancy between the expected and observed solar neutrino flux introduces, however, a certain degree of uncertainty in the conclusions of Ref. 17. Thus, the present limits in Fig. 4, although not as restrictive as in other ranges, represent a useful addition.

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