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Astrophysical Bounds on Neutrino Radiative Lifetimes Revisited*+

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We critically re-examine the limits on neutrino radiative lifetimes that can be obtained from the absence of solar γ-rays, from the observed flux of γ-rays from the galactic center, and from the observed background spectrum of diffuse, isotropic x- and γ-rays. We find $\tau_{\nu_e}/m_{\nu_e} > 7{\cdot}10^9\,\mathrm{sec/eV}$ from the solar neutrino flux and a similar result from the galactic center. The dependence of this result on the energy of the decay photon and on neutrino masses and mixing parameters is discussed. The diffuse cosmic neutrino flux from hydrogen burning stars yields $\tau_{\nu_e}/m_{\nu_e} > 10^{12}\,\mathrm{sec/eV}$ and the same limit holds from the diffuse flux of supernova neutrinos. In the latter case the bound also applies to μ-neutrinos. The limit from the diffuse flux from degenerate dwarf formation is more restrictive by about an order of magnitude. — The original application of several of these arguments by R. COWSIK yields bounds which are 2–4 orders of magnitude more restrictive, partly due to his misrepresentation of observational data, partly due to his choice of cosmological parameters which we believe to be inconsistently optimistic. — In parts of this work we follow a recent discussion by P. VOGEL on the decay of reactor neutrinos.

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1. Introduction

If neutrinos were massive they could be unstable and would presumably exhibit a radiative decay mode of the type

$$\nu \longrightarrow \nu' + \gamma . \tag{1}$$

ν' is some neutral fermion, possibly some lighter neutrino species. The standard model of electroweak interactions provides detailed predictions on the radiative decay widths in terms of the neutrino masses and mixing-parameters for the case where the flavor-eigenstates ν_e, ν_μ, ν_τ, ... do not coincide with the mass-eigenstates ν_1, ν_2, ν_3, ... [1, 2].

Independently of such predictions and without specifying the identity of ν' the conversion rate of various neutrino fluxes into radiation through (1) have been limited from existing experimental or observational data. After the relevant neutrino sources these limits can be subdivided into the classes of laboratory, astrophysical and cosmological limits. The latter ones are based on the thermal neutrino flux expected from a hot early stage of the universe. Due to the large flux and long decay-times available in this context, mass-dependent results of up to $\tau_\nu \gtrsim 10^{25}$ sec have been obtained,[1] far exceeding the age of the universe of $4\text{–}6 \cdot 10^{17}$ sec. The cosmic neutrino flux, however, is known only from calculations in the framework of standard Big Bang cosmology and is, at present, not accessible to direct observation. Therefore limits based on more directly established neutrino fluxes still deserve attention in spite of their less restrictive values.

The standard[2] laboratory limit on the electron (anti-)neutrino[3] radiative lifetime, $\tau_{\nu_e}/m_{\nu_e} > 300\,\text{sec/eV}$, was derived by REINES et al. [4] from the absence of single photon counts in a detector near a fission reactor which produces an anti-neutrino flux of known intensity and spectral distribution. According to a recent criticism by VOGEL [5] this number should be reduced by an order of magnitude. Using data from the Gösgen-reactor and employing a much more detailed analysis, VOGEL derives a new limit of $\tau_{\nu_e}/m_{\nu_e} > 1\,\text{sec/eV}$. The standard laboratory limit

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[1] For a brief review of these limits and detailed references see the appendix to this note.

[2] This refers to the 1984 edition of the *Review of Particle Properties* [3].

[3] If neutrinos have nonzero masses, as we assume here, the distinction between Dirac- and Majorana-neutrinos becomes important. In the latter case the term "anti-neutrino" is understood to refer to the same particle as the term "neutrino".

on the ν_μ radiative lifetime [6] is $\tau_{\nu_\mu}/m_{\nu_\mu} > 0.11\,\mathrm{sec/eV}$, and no such bounds have been reported for ν_τ.

As was first pointed out by COWSIK [7], these numbers may be heavily improved upon by using astrophysical neutrino fluxes such as that from the sun, from white dwarf formation or from supernova events. The conversion of these neutrino fluxes during astronomical decay times into radiation is limited by the observational bounds [8, 9] on the solar x- and γ-ray spectrum and by the measured background spectrum [10] of diffuse, isotropic x- and γ-rays.

COWSIK's order-of-magnitude estimates, however, result in bounds which are too restrictive by 2–4 orders of magnitude, as we shall argue in detail in the present discussion. Therefore we reconsider the bounds on the neutrino radiative decay widths from astrophysical data in the spirit of COWSIK's original work [7]. In our discussion of the solar bounds we follow in some detail a recent discussion [5] by VOGEL on the decay of reactor neutrinos.

In addition to the astrophysical neutrino sources suggested in COWSIK's original treatment—the sun and the cosmic diffuse neutrino flux from stellar collapse—we consider the hydrogen burning stars in the central bulge of the Milky Way and speculate on what could be learned from the neutrino flux emerging from collapsing stars in the Virgo cluster of galaxies. These quasi discrete objects are reasonable neutrino sources in comparison with the sun because the flux of decay photons suffers a dilution which is proportional only to $(\mathrm{distance})^{-1}$. In the case of the galactic center and the Virgo cluster the increased luminosity in comparison with the sun makes up for this loss in photon flux which is so mild because neutrinos from a more distant source have more time available for their decay. Furthermore the cosmic diffuse neutrino flux from hydrogen burning stars is considered as a source for decay photons which could contribute to the diffuse cosmic flux of x- and γ-rays.

The sun and other hydrogen burning stars produce MeV–neutrinos in charged current reactions which amout to $p + e^- \rightarrow n + \nu_e$. Therefore the relevant bounds only apply to electron neutrinos and, with modifications, to hypothetical heavy neutrino admixtures to this species. Thermal neutrinos from supernovas and degenerate dwarf formation are largely produced in neutral current reactions and therefore the limits derived from such sources apply to all neutrino flavors light enough to be produced in the respective heat baths.

2. Solar Bounds on Neutrino Radiative Lifetimes

2.1 Order of magnitude estimate for ν_e

The sun in believed to produce energy by the fusion of hydrogen into helium. This process corresponds to the effective reaction

$$4p + 2e^- \longrightarrow \alpha + 2\nu_e + 26.73\,\text{MeV}, \tag{2}$$

which proceeds through various reaction chains and cycles. Then the total solar neutrino flux at earth of about $6.6 \cdot 10^{10}\,\text{cm}^{-2}\,\text{sec}^{-1}$ is easily calculated from the solar electromagnetic luminosity of $2.4 \cdot 10^{39}$ MeV/sec and from the distance to the sun of $1.50 \cdot 10^{13}$ cm by observing that only a few percent of the above reaction energy is carried away by the neutrinos.

The relative contributions of the four reaction chains *PP I–III* and *CNO* depend on details of the condition in the solar interior; a standard calculation [11] results in the solar neutrino spectrum as displayed in Fig. 1, curve (a).[4] The high energy part—above approximately 1 MeV—of this spectrum has been measured in a pioneering experiment [12], yielding a flux which is approximately a factor 3–4 below the calculated amplitude [11]. In spite of this "solar neutrino puzzle" these measurements render the solar neutrino flux the ***only observed*** and hence the ***only directly established extra-terrestrial neutrino flux.***[5] The dominant low energy part of this spectrum will probably be measured within the next few years with a Gallium neutrino detector [15]. The anticipated confirmation of the calculated flux would promote the solar neutrino radiative lifetime limits from an almost to an exactly equal footing as a laboratory result.

While the calculation of the high energy part of the solar neutrino spectrum involves fine points of solar modelling and ill-known nuclear reaction parameters [11], it appears safe to assume that approximately 90 % or more of all solar neutrinos are produced in the pp-reaction

$$p + p \longrightarrow d + e^+ + \nu_e + 420\,\text{keV} \tag{3}$$

[4] A similar plot of the calculated neutrino spectrum in the standard reference [11], on which our spectrum is based, is inconsistent with their calculated neutrino flux: The amplitude of the high energy part (from the decay of ^{8}B) is plotted almost an order of magnitude too large.

[5] We note, however, that LANDE et al. [13] have speculated that an anomalous series of pulses in their $\bar{\nu}_e$-detector was due to a pulse from a collapsing star in the center of the galaxy. Very recently DAR [14] has shown that a background of neutrino induced events in the KAMIOKANDE proton decay detector is compatible with the diffuse neutrino flux expected from collapsing stars in the whole universe.

with a spectrum of an allowed β-decay and an average neutrino energy of 265 keV. Then we assume for now that neutrinos can decay via reaction (1) with (partial) proper lifetime τ_{ν_e}, and that the energy of 90 % of all solar neutrinos is *exactly* $E_{\nu_e} = 265\,\mathrm{keV}$. During the time of flight from sun to earth, $t_\odot = 499\,\mathrm{sec}$, the exponential decay law can be linearly approximated since we know from VOGEL's result [5] that the laboratory decay time of neutrinos with the above energy exceeds $2.7{\cdot}10^5$ sec. This number also guarantees that the neutrinos do not decay with a substantial probability before reaching the solar surface. Hence a fraction[6] $(m_{\nu_e}/E_{\nu_e}) \cdot (t_\odot/\tau_{\nu_e})$ of the solar neutrino flux at this energy, $F_{\nu_e} = 0.90 \cdot 6.6{\cdot}10^{10}\,\mathrm{cm}^{-2}\,\mathrm{sec}^{-1} = 5.9{\cdot}10^{10}\,\mathrm{cm}^{-2}\,\mathrm{sec}^{-1}$, is converted into photons, yielding a *total* photon flux at earth of $F_\gamma = (m_{\nu_e}/E_{\nu_e}) \cdot (t_\odot/\tau_{\nu_e}) \cdot F_{\nu_e}$. This flux exhibits a nontrivial distribution in photon-frequency ω. If we assume that the particle ν' is massless, that the neutrino decay is isotropic in its own rest frame, and that $m_{\nu_e} \ll E_{\nu_e}$, the photon spectrum will be constant between 0 and E_{ν_e} and zero elsewhere. Hence we obtain for the *differential* photon flux at earth the "boxspectrum"

$$j_\gamma(\omega) = \begin{cases} \dfrac{m_{\nu_e}}{\tau_{\nu_e}} \cdot \dfrac{t_\odot}{E_{\nu_e}^2} \cdot F_{\nu_e} & \text{for} \quad 0 \leq \omega \leq E_{\nu_e}, \\ 0 & \text{for} \quad \omega > E_{\nu_e}, \end{cases} \tag{4}$$

where the r.h.s. is $(m_{\nu_e}/\tau_{\nu_e}) \cdot (\mathrm{sec/eV}) \cdot 4.2{\cdot}10^5\,\mathrm{cm}^{-2}\,\mathrm{sec}^{-1}\,\mathrm{keV}^{-1}$ in the case at hand.

The photon flux $j_\gamma(\omega)$ expected from neutrino decay must lie below the observational bounds or measurements of the solar x- and γ-radiation. In Fig. 1, curve (c) we show the upper limit photon spectrum of the quiet sun from balloon flight measurements [8,9] in the range $20\,\mathrm{keV} \lesssim \omega \lesssim 10\,\mathrm{MeV}$. For energies below approximately 10 keV the emission is dominated by an intense bremsstrahlung spectrum—Fig. 1, curve (d)—of the hot solar corona ($T \approx 4.5{\cdot}10^6\,\mathrm{K}$) which is observationally established from rocket flight measurements [16].

In the interval $0 \leq \omega \leq 265\,\mathrm{keV}$ the lowest limit is $0.6{\cdot}10^{-4}\,\mathrm{cm}^{-2}\,\mathrm{sec}^{-1}\,\mathrm{keV}^{-1}$. Hence we find a lower bound of

$$\tau_{\nu_e}/m_{\nu_e} > 7{\cdot}10^9\,\mathrm{sec/eV}\,. \tag{5}$$

COWSIK [7] erroneously finds instead $\tau_{\nu_e}/m_{\nu_e} > 2{\cdot}10^{12}\,\mathrm{sec/eV}$, 2.5 orders of magnitude above our result, because he quotes from ref. [9] an *integrated* flux limit

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[6] We use natural units with $\hbar = c = 1$ throughout this paper.

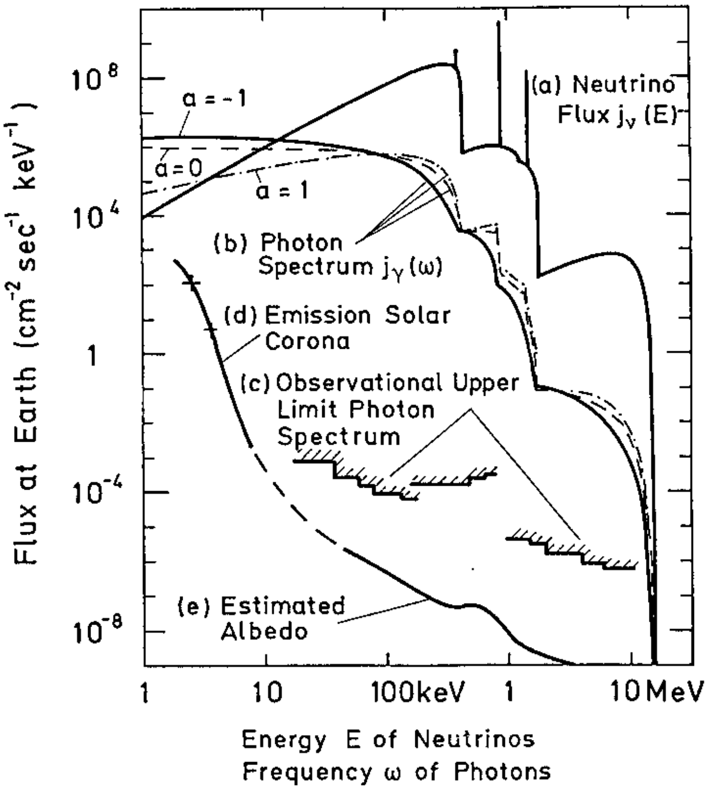

Fig. 1 **(a)** Predicted solar neutrino flux $j_{\nu_e}(E)$ according to [11]. The leftmost "bump" stems from the pp-reaction, $p + p \rightarrow d + e^+ + \nu_e$, and contributes $6.1 \cdot 10^{10}\,cm^{-2}sec^{-1}$ to the total flux of $6.6 \cdot 10^{10}\,cm^{-2}sec^{-1}$ at the earth. The three "spikes" correspond to electron-capture reactions. Their width is dominated by the thermal motion of the sources in the solar interior. For convenience we have normalized them such as if they had rectangular shape and a width of exactly 1 keV. — **(b)** Spectrum of photons $j_\gamma(\omega)$ from the decay of the solar neutrino spectrum (a) through the process $\nu_e \rightarrow \nu' + \gamma$. ν' is assumed to be massless, corresponding to $r = 1$, and $\tau_{\nu_e}/m_{\nu_e} = 1\,sec/eV$ is used. — **(c)** Upper limit on the x- and γ-radiation of the quiet sun according to [8, 9]. — **(d)** Emission of the hot solar corona ($T = 4.5 \cdot 10^6\,K$) according to [9]. The data-points "+" are taken from [16]. — **(e)** Estimated solar albedo according to [9].

of $10^{-4}\,cm^{-2}\,sec^{-1}$ although this reference gives the above ***differential*** flux limit. Therefore he misses a factor $E_{\nu_e} = 265\,keV$ in the denominator of his result, accounting for the discrepancy between his and our limit.

For 265 keV neutrinos our result translates into a ***laboratory*** lifetime limit of $\tau_{\nu_e} > 2 \cdot 10^{15}\,sec.$[7] If there are other decay channels open such as $\nu \rightarrow 2\nu' + \bar{\nu}'$, the neutrino flux must not be reduced too much by these decays or our analysis for the radiative decay mode will be invalidated. If the branching ratio of the radiative decay exceeds about 10^{-12}, the total laboratory lifetime will still exceed the time of flight from sun to earth and our analysis remains valid.

[7] COWSIK's [7] corresponding *laboratory* lifetime limit of $\tau_{\nu_e} > 5 \cdot 10^{17}\,sec$ has been quoted by DOLGOV & ZELDOVICH in their review "Cosmology and Elementary Particles" [17] such as if it were a *proper* lifetime limit, thereby further obscuring the neutrino radiative lifetime limits that can be obtained from the sun.

2.2 Detailed analysis

Now we relax most of the simplifying assumptions of our above order-of-magnitude analysis and only maintain the notion of an electron neutrino which is essentially a mass-eigenstate. If it decays in a process of type (1) the energy of the decay photon in the neutrino rest frame is

$$\omega = r \cdot m_{\nu_e}/2\,, \quad 0 < r \leq 1\,, \tag{6}$$

where m_{ν_e} is the mass of the electron neutrino. The dimensionless parameter r is the ratio between the actual photon energy ω and its possible maximum value, $m_{\nu_e}/2$, which is assumed when the particle ν' is massless. In terms of the neutrino masses we find

$$r = 1 - \left(\frac{m_{\nu'}}{m_\nu}\right)^2. \tag{7}$$

Limits on neutrino radiative lifetimes have generally been derived under the assumption that the mass of ν' can be neglected. Such a constraint, however, is not necessary and we prefer to derive the lifetime limit as a function of the parameter r.

Following VOGEL [5], we write the most general form of the angular distribution of the decay photons in the ν_e rest frame for the case where ν' carries spin $\frac{1}{2}$ as

$$\frac{dN}{d\cos\theta} = \frac{1}{2}(1 + a\cos\theta), \quad |a| \leq 1\,. \tag{8}$$

θ is the angle between the directions of motion of ν_e and γ. It is assumed that the polarization vector of ν_e is (anti-)parallel to its momentum. For Majorana-neutrinos the decay would be isotropic, independently of their polarization, and thus $a = 0$. For left-handed (helicity-minus) Dirac neutrinos and a massless ν' one has $a = -1$, see ref. [1, 2], meaning that the photons are preferably emitted "backward". The exactly same is true for the corresponding helicity-plus Dirac anti-neutrino states as is required by CP-invariance and, if final state interactions can be neglected, CPT-invariance.

From the calculated solar neutrino flux $j_{\nu_e}(E)$ at the earth—see ref. [11] and Fig. 1, curve (a)—and using (8) one can then derive the expected differential photon flux from neutrino decay[8] for the case $m_{\nu_e} \ll E$

$$j_\gamma(\omega) = \frac{1}{r} \cdot \frac{m_{\nu_e}}{\tau_{\nu_e}} \cdot t_\odot \int_{\omega/r}^{\infty} \left(1 - a + 2a\frac{\omega/r}{E}\right) \frac{j_{\nu_e}(E)}{E^2}\, dE\,. \tag{9}$$

In Fig. 1, curve (b) we display this function for $r = 1$, $\tau_{\nu_e}/m_{\nu_e} = 1\,\mathrm{sec/eV}$ and $a = 0, \pm 1$. For $r = 1$, $a = 0$ and $j_{\nu_e}(E) = F_{\nu_e} \cdot \delta(265\,\mathrm{keV} - E)$ this reduces, of course, to the "boxspectrum" (4).

The structure of this function is such that the spectrum for other values of r can be obtained—in a doubly logarithmic representation such as Fig. 1—by simply shifting it "to the left" by the amount $|\log_{10} r|$ and "upward" by the same amount, the shape itself remaining unchanged. For $r = 10^{-1}$, as an example, it is shifted to the left and upward by one decade each.

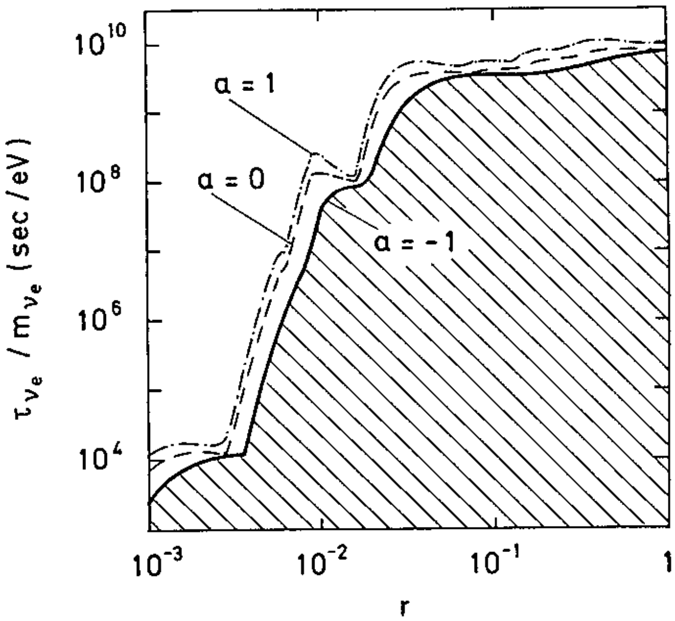

Fig. 2 Possible values for the electron neutrino's lifetime-over-mass-ratio τ_{ν_e}/m_{ν_e} and for the energy fraction $r = \omega/(m_{\nu_e}/2)$ which is carried away by the photon in the reaction $\nu_e \to \nu' + \gamma$ are restricted to the white area in this plot. Curve $a = -1$ yields the most conservative limits.

The requirement that the photon spectrum must lie below the observational bounds excludes the shaded area of $(r, \tau_{\nu_e}/m_{\nu_e})$-values in Fig. 2. Since $a = -1$ is the worst case in the sense that the photon spectrum is squeezed furthest toward low frequencies, the corresponding curve in Fig. 2 yields the most conservative limit on τ_{ν_e}/m_{ν_e}. Then we obtain for $r = 1$ the same limit (5) as from our above

[8] There is an apparent misprint concerning the sign of a in the corresponding eq. (6) of ref. [5].

order-of-magnitude calculation. For $0.03 \leq r \leq 1$ our limit remains above the 10^9 sec/eV-level, while it drops quite fast for lower r-values. In the range above 0.03 we consider the limit firm, because it depends only on the low-energy part of the solar neutrino spectrum. This range of r-values corresponds from (7) to the case where the mass difference between ν and ν' exceeds about 1.5 %.

The observational upper limit photon flux is several orders of magnitude above the expected solar albedo—Fig. 1, curve (e)—leaving much room for an observational improvement of our limits. The actual solar albedo could be much lower than this estimate, because the cosmic ray flux hitting the surface of the sun could be substantially reduced from the galactic average due to shielding effects by the solar magnetic field [9]. In principle our limits could be further improved by a factor of about 25 if one could achieve an angular resolution of a few minutes of arc, because neutrino decay photons would appear to come from the center of the solar disc. This should be so because the solar neutrinos are produced in the center of the sun and the low mass limit for electron neutrinos guarantees an ultrarelativistic γ-factor which suppresses strong angular deviations between ν_e and γ in the observer's frame.

2.3 Light, subdominant ν_e-admixtures

In contrast to our above assumption the electron neutrino could be a mixture of mass eigenstates ν_i with masses $m_1 \leq m_2 \leq \ldots \leq m_n$,

$$\nu_e = \sum_{i=1}^{n} U_{ei}\, \nu_i\,, \tag{10}$$

where ν_e is understood to be the state that couples to electrons in charged current weak interactions. A decay of ν_e would then really correspond to a decay of its heavy neutrino admixtures into radiation and lighter neutrino states. The standard model of weak interactions provides detailed predictions [1] on the relevant decay widths for this case. We note that there is no a-priori-reason to identify ν_e mainly with ν_1, ν_μ mainly with ν_2 etc.; ν_e could be essentially the heaviest neutrino species. Our previous results then refer to the dominant ν_e-admixture which could be any one of $\nu_1, \ldots, \nu_n$ and which is known to be lighter than about 46 eV.

Turning to the subdominant admixtures we note that the differential neutrino flux for *any* neutrino component which is so light that its production in the sun is not inhibited by threshold effects ($m_i \lesssim 60\,\text{keV}$) is obtained from $j_{\nu_e}(E)$ by multiplication with $|U_{ei}|^2$. The spectrum of decay photons is obtained by substituting τ_i/m_i for τ_{ν_e}/m_{ν_e} in our previous discussion. Then our previous limits translate to the present case through the substitution

$$\frac{\tau_{\nu_e}}{m_{\nu_e}} \longrightarrow |U_{ei}|^{-2} \cdot \frac{\tau_i}{m_i}. \tag{11}$$

If $r \approx 1$ the previous limit (5) together with the above substitution yields the bound

$$\tau_i > 7{\cdot}10^9\,\text{sec} \cdot \frac{m_i}{1\,\text{eV}} \cdot |U_{ei}|^2 \quad \text{for} \quad m_i \lesssim 60\,\text{keV}. \tag{12}$$

For all components ν_i for which $r \gtrsim 0.03$ and $m_i \lesssim 60\,\text{keV}$ this means

$$\left(\sum_{i=2}^{n} |U_{ei}|^2 \frac{m_i}{\tau_i}\right)^{-1} \gtrsim 10^9\,\text{sec/eV}. \tag{13}$$

2.4 Heavy, subdominant ν_e-admixtures

Now we consider the case where $10\,\text{keV} \leq m_i \leq 1\,\text{MeV}$ for some mass eigenstate ν_i. We note that the standard mass limit $m_{\nu_\mu} < 0.50\,\text{MeV}$ guarantees that $m_2 < 1\,\text{MeV}$. Therefore at least ν_2 could lie in this range. If m_i were much in excess of 1 MeV the decay $\nu_i \to \nu_1 + e^- + e^+$ would probably dominate over the radiative decay-mode. Limits for this case have been obtained by TOUSSAINT & WILCZEK from the low energy interplanetary positron density [18]. Since $m_1 < 46\,\text{eV}$ it is safe to assume $r \approx 1$, i.e. to treat ν_1 as a massless particle in comparison with ν_i for the present considerations.

The differential neutrino flux is now

$$j_{\nu_i}(E) = \begin{cases} 0 & \text{for} \quad 0 \leq E \leq m_i, \\ |U_{ei}|^2 \cdot \beta(E) \cdot j_{\nu_e}(E) & \text{for} \quad m_i < E \lesssim 15\,\text{MeV}, \end{cases} \tag{14}$$

where $\beta(E) = \sqrt{1-(m_i/E)^2}$ is the velocity of a neutrino ν_i of energy E. Hence this spectrum depends now sensitively on the neutrino mass due to its nonrelativistic phasespace factor in the production rate.

The present case is further complicated by the fact that the component ν_i is not relativistic anymore for the whole range of relevant energies. Therefore the decay photons do not appear to come from the solar center anymore but rather exhibit an angular distribution which is symmetric about the axis sun–observer. At an angle ϑ away from the sun and at a photon frequency ω we find that the differential photon flux at the earth vanishes for all neutrino energies unless (ω, ϑ) lie within one of the following parameter ranges:

$$\begin{array}{llll} \text{(i)} & 0 < \omega < \frac{m_i}{2} & \text{and} & 0 \leq \vartheta \leq \pi \\ \text{(ii)} & \frac{m_i}{2} \leq \omega < \infty & \text{and} & 0 \leq \vartheta \leq \arcsin\left(\frac{m_i}{2\omega}\right). \end{array} \tag{15}$$

We note that in case (ii) the absolute maximum value for ϑ is $\frac{\pi}{2}$. This means that decay photons with energies below $m_i/2$ can be observed from all directions, while photons with higher energies are more and more constrained to directions near the sun and never occur further away than $\vartheta = 90^\circ$.

Then we find for the doubly differential flux $j_\gamma(\omega) = dF_\gamma/d\omega d\Omega$ the following result where we have used the velocity β as a variable to integrate over the neutrino source spectrum:

$$j_\gamma(\omega, \vartheta) = |U_{ei}|^2 \cdot \int\limits_{\beta_{\min}}^{\beta_{\max}} d\beta \cdot j_{\nu_e}\left(\frac{m_i}{\sqrt{1-\beta^2}}\right) \cdot \frac{\beta + a \cdot \left(\frac{2\omega}{m_i}\sqrt{1-\beta^2} - 1\right)}{1-\beta^2} \cdot \frac{x(\omega,\beta) \cdot e^{-x(\omega,\beta)\sin\vartheta}}{2\pi \sin\vartheta} \tag{16}$$

where

$$x(\omega, \beta) = \frac{t_\odot}{\tau_i} \cdot \frac{\omega}{\sqrt{\frac{m_i \omega}{\sqrt{1-\beta^2}} - \frac{m_i^2}{4} - \omega^2}}$$

and

$$\beta_{\min} = \begin{cases} \dfrac{m_i^2 - 4\omega^2}{m_i^2 + 4\omega^2} & \text{for range (i),} \\[2ex] \dfrac{4\omega^2 \cos\vartheta - m_i\sqrt{m_i^2 - 4\omega^2 \sin^2\vartheta}}{m_i^2 + 4\omega^2\cos^2\vartheta} & \text{for range (ii),} \end{cases}$$

$$\beta_{\max} = \frac{4\omega^2 \cos\vartheta + m_i\sqrt{m_i^2 - 4\omega^2 \sin^2\vartheta}}{m_i^2 + 4\omega^2\cos^2\vartheta}.$$

We note that the parameter a which governs the angular distribution of the photons in the neutrino rest frame through eq. (8) and hence the energy distribution in the laboratory frame does not necessarily take on the same constant value for all velocities β, because nonrelativistic neutrinos emerging from a weak interaction process would be polarized proportional to β. In order to derive conservative bounds we have always used such a function $a(\beta)$ that the integrand in (16) is minimal.

Now there are two ways to constrain the neutrino input parameters $|U_{ei}|^2$, m_i and τ_i from observational data. The doubly differential flux $j_\gamma(\omega, \vartheta)$ for all allowed values ω and ϑ must lie below the measured background [10] of diffuse, isotropic x- and γ-rays. It turns out that for certain values of the parameters $(|U_{ei}|^2, m_i, \tau_i)$ the calculated flux $j_\gamma(\omega, \vartheta)$ is quite isotropic aside from angles near $\vartheta = 0^\circ$ and $\vartheta = 180^\circ$. Therefore solar neutrino decay may contribute to the diffuse photon background.

The constraints such obtained, however, are slightly weaker than can be derived from the observational bounds on the direct solar x- and γ-radiation which we have used before. These bounds were obtained by balloon borne photon-detectors which looked toward the sun, registering photons within a finite forward aperture. Then the same detector looked into the exactly opposite direction for some time, and the two counting rates were subtracted from each other to derive the observational upper limit flux.

If the forward aperture covers an opening angle $\Omega_{\rm ap}$ and if axisymmetry is assumed, this corresponds to a maximum angle $\vartheta_{\rm max} = \arccos(1 - \Omega_{\rm ap}/2\pi)$ from which photons can be registered. In the two relevant cases these values were

$\Omega_{\rm ap}$	$\vartheta_{\rm max}$	Photon-energy	Ref.	
≈ 0.15 ster	$\approx 12^\circ$	18–185 keV	PETERSON et al. [9]	(17)
≈ 1 ster	$\approx 33^\circ$	163–774 keV	FROST et al. [8].	

We follow the experimental procedure and calculate the effective photon flux

$$j_{\rm eff}(\omega) = 2\pi \left(\int_0^{\vartheta_{\rm max}} - \int_{(\pi-\vartheta_{\rm max})}^{\pi} \right) j_\gamma(\omega, \vartheta) \sin\vartheta \, d\vartheta \,. \tag{18}$$

The implicit assumption of isotropic response of the detector within the forward aperture is, of course, not strictly justified.

This effective flux again must lie below the observational bounds—Fig. 1, curve (c)—for each choice of $(|U_{ei}|^2, m_i, \tau_i)$. This requirement translates into lower

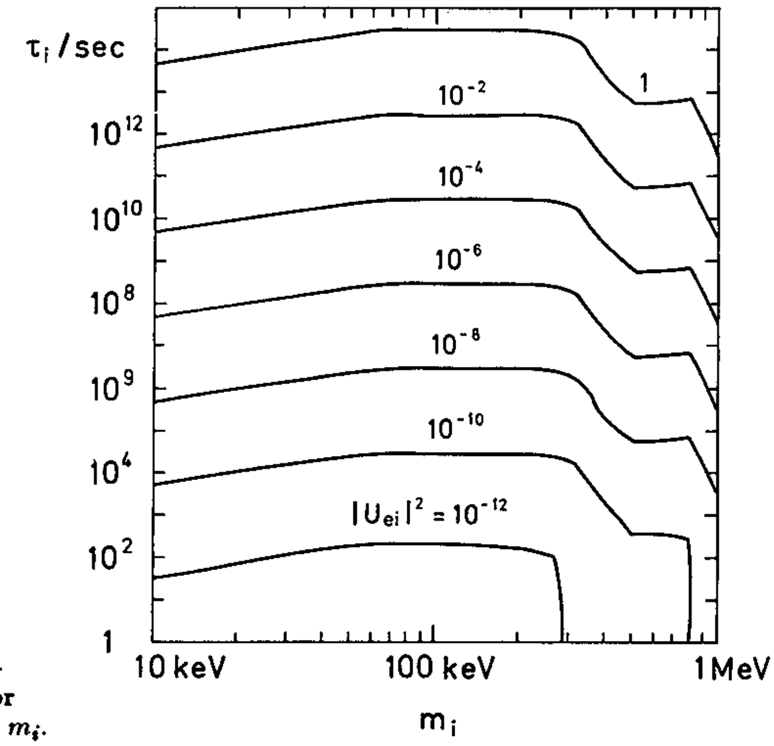

Fig. 3 Each curve gives for the indicated value of $|U_{ei}|^2$ the lower limit for the lifetime τ_i as function of the mass m_i.

bounds for the lifetime τ_i which are graphically represented in Fig. 3. There the lower limits are plotted as a function of m_i for several values of $|U_{ei}|^2$. We note that these curves vary only in amplitude, not in shape, for $\tau_i \gtrsim 10^3$ sec. For shorter lifetimes the exponential form of the decay law becomes important, accounting for the distortion of the curves in the "lower right corner" of Fig. 3. For lifetimes much shorter than 1 sec no limits can be derived, because then most neutrinos would decay before reaching the solar surface.

The present results can be expressed for $\tau \gtrsim 10^3$ sec approximately as

$$\tau_i > 5{\cdot}10^{13}\,\mathrm{sec} \cdot \frac{m_i}{10\,\mathrm{keV}} \cdot |U_{ei}|^2 \qquad \text{for} \qquad 10\,\mathrm{keV} \leq m_i \lesssim 60\,\mathrm{keV} \tag{19}$$

and

$$\tau_i > 2{\cdot}10^{14}\,\mathrm{sec} \cdot |U_{ei}|^2 \qquad \text{for} \qquad 60\,\mathrm{keV} \lesssim m_i \lesssim 300\,\mathrm{keV}. \tag{20}$$

The first of these results is practically identical with our previous limit (12). We note that the present results depend on a higher energy region of the solar neutrino spectrum than previously used. They are still independent of the neutrinos expected from ^8B decay.

3. Limits from Other Discrete Sources

3.1 The central bulge of the Milky Way

Not only the sun, but all hydrogen burning stars gain the energy for their luminosity from reaction (2) and hence produce MeV–neutrinos at a similar rate as the sun. It is hopeless, however, to consider any other isolated star, because the expected flux of decay photons from a distant neutrino source scales with $(\text{distance})^{-1}$. This is different from the $(\text{distance})^{-2}$–scaling of the neutrino flux itself, because for large neutrino lifetimes the fraction of this flux which is converted into photons scales with $(\text{distance})^{+1}$.

One may consider, however, such a large association of stars as the central bulge of the Milky Way. Within a radius of 2.5 kpc it contains a luminosity of about $2{\cdot}10^{10}\,L_\odot$ and with its distance of $(8.7 \pm 0.6)\,\text{kpc}$ [19] it is about $2{\cdot}10^{9}$ times farther away from us than the sun.[9] Then one expects about an order of magnitude more decay photons at the earth from there than from the sun.

The x- and γ-ray flux from the galactic center has been measured in various experiments [20], the measured fluxes being similar in magnitude to the upper limit solar fluxes. Since the above values for distance and size of the central bulge of the Milky Way imply an angular size of about 30° we are interested in results obtained with a detector with a large forward aperture so that it would see a large fraction of the central bulge. According to ref. [20] such data are available from the HEAO–1 mission of September 1977 where the detector had a 16° field of view in the range 100 keV to 1 MeV and 43° between 1 and 10 MeV. Then a conservative upper bound to the emission of the central bulge in the range 100 keV to 10 MeV is

$$j_\gamma(\omega) = b \cdot \omega^{-2} = 2{\cdot}10^{-5}\,\text{cm}^{-2}\,\text{sec}^{-1}\,\text{keV}^{-1} \cdot \left(\frac{1\,\text{MeV}}{\omega}\right)^2 \qquad (21)$$

where we have allowed for at least a factor of 2 for the reduction of the measured flux from the true value due to the reduced detection efficiency of a source off the aperture axis.

Now we require that the expected photon spectrum (9) must lie below the observational upper bound (21). $j_{\nu_e}(E)$ now means the neutrino flux from the galactic center and $t_\odot$ must be replaced by the time of flight from the galactic

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[9] $1\,\text{kpc} = 3.08{\cdot}10^{21}\,\text{cm}$; $1\,L_\odot = 2.40{\cdot}10^{39}\,\text{MeV/sec}$.

center to the earth, $t_\otimes \approx 9{\cdot}10^{11}$ sec. Turning to the worst case with $a = -1$ and using $r = 1$ we then find the bound

$$\frac{\tau_{\nu_e}}{m_{\nu_e}} > \frac{t_\otimes}{b} \cdot \max_{0 \le \omega \le E_{\max}} 2 \int_{\omega}^{E_{\max}} \left(1 - \frac{\omega}{E}\right)\left(\frac{\omega}{E}\right)^2 j_{\nu_e}(E)\, dE \,. \tag{22}$$

When the neutrinos have definite energy E_0, $j_{\nu_e}(E) = F_{\nu_e} \cdot \delta(E_0 - E)$, this result simply reads

$$\frac{\tau_{\nu_e}}{m_{\nu_e}} > \frac{t_\otimes}{b} \cdot \frac{8}{27} \cdot F_{\nu_e} \,. \tag{23}$$

It is independent of E_0 due to the ω^{-2}–dependence of the observational bounds.

In order to translate the bolometric luminosity of the galactic center into a neutrino luminosity we note that the dominant stellar energy resource is nuclear binding energy. Gravitational binding energy of compact stellar objects rivals this resource only in the case of the formation of neutron stars or black holes. Such events, however, are rare and radiate away most of the liberated binding energy as energetic neutrinos which we do not consider as a neutrino source at this point of our discussion.

Therefore the bolometric luminosity of galaxies can be mainly attributed to the liberation of nuclear binding energy. Nuclear synthesis proceeds first from hydrogen to helium where 2 neutrinos are emitted per synthesized α-particle. In more advanced stellar burning stages essentially α-particles are combined to form heavier nuclei such as ^{12}C, ^{16}O, ^{20}Ne, etc.; no neutrinos are produced in such reactions. However, about 6.7 MeV of the approximately 8.8 MeV of nuclear binding energy per nucleon that can be maximally gained from nuclear fusion is liberated during hydrogen burning. Therefore it is safe to assume that at least 75 % of the electromagnetic energy output of a galaxy can be attributed to hydrogen burning.

Hydrogen burning is believed to proceed mainly through the reaction chains *PP I–III* and *CNO* [21] which are of different relative importance at different stellar temperatures, densities and chemical compositions. Each of them produces a characteristic neutrino spectrum. In order to avoid assumptions about the relative average occurrences of these chains in the galaxy we choose them such that our results will be most conservative. This means that we vary the relative rates until that the r.h.s. of eq. (22) becomes minimal while keeping constant the total bolometric luminosity of $0.75 \cdot 2{\cdot}10^{10}\,L_\odot$ due to hydrogen burning.

We find this to be the case when the *PPIII* chain and *CNO* cycle each contribute half to that part of the galactic energy production which does not go into neutrinos. We note that for *PPIII* on average 28 % of the energy liberated in reaction (2) is carried away by neutrinos, this number being 6 % for *CNO*, 2 % for *PPI* and 4 % for *PPII*.

In this somewhat unrealistic but conservative case we find the bound

$$\frac{\tau_{\nu_e}}{m_{\nu_e}} > 1.6\cdot 10^9\,\mathrm{sec/eV}\,, \tag{24}$$

slightly less restrictive than the solar limit. This result is still interesting because it is derived from a completely different set of observational data. It could be improved at the most by a factor of 2 by a detailed discussion of the average rate of occurrence of the various reaction chains, i.e. by a detailed calculation of the galactic neutrino spectrum due to hydrogen burning. It could be improved observationally if the x- and γ-ray flux from the galactic center could be resolved in its angular structure and then attributed to a discrete source, possibly the galactic nucleus [20].

3.2 Speculations on supernovas in the Virgo cluster

Beyond the center of our galaxy one may consider other discrete sources at larger distances if they are appropriately more luminous. The only obvious such candidate is the Virgo cluster of galaxies at a distance of about 15.7 Mpc [22], about 1800 times farther away than the center of the Milky Way. The luminosity of its estimated number of 2500 galaxies [23] could approximately compensate for the $(\mathrm{distance})^{-1}$-loss in decay photons.

As in the case of the galactic center, the solar bounds could be only improved by much lower observational limits to the Virgo clusters's x- and γ-ray flux. Alternatively, an improvement could result from the consideration of other neutrino sources in the Virgo cluster but hydrogen burning stars, e.g., collapsing stars. As we shall presently discuss, the time-averaged neutrino flux from neutron star formation is similar in magnitude to the flux from hydrogen burning stars, and the same would apply to the flux of decay photons. Its pulsed nature, however, would momentarily raise the observable flux to much larger values. In the case of the galactic center one cannot take advantage of this fact because of the rare occurrence of neutron star formation which is believed to take place only a few times per century in a galaxy. In the Virgo cluster, however, one expects such events once every few weeks. This time scale encourages us to speculate in this section on the neutrino lifetime limits that could be achieved from the observational absence of x- and γ-ray bursts preceding optical supernova events.

Neutron star formation is believed to be accompanied by type II supernova explosions, although it is not known whether they can also be formed occasionally in type I events and whether a quiet birth is possible [24]. If a neutron star of 1 $M_\odot$ is formed one expects a pulse of about $6\cdot 10^{56}$ electron neutrinos from the "deleptonization of matter", $p + e^- \to n + \nu_e$. Furthermore there will be thermal emission of neutrinos of all flavors. Considering numbers from specific model calculations [25, 26] we adopt the value

$$N_f = x_f \cdot 10^{58} \tag{25}$$

for the total number of neutrinos plus antineutrinos per flavor f emitted in a supernova event. x_f is a flavor dependent "fudge factor" of order unity. These neutrinos have average energies of 10 to 20 MeV and a width in the energy distribution of similar magnitude. The duration of the neutrino pulse is between a few hundred milliseconds and a few seconds.

Due to the uncertainty about the correlation of neutron star formation with type II supernovas it is difficult to estimate precisely the average rate of such collapse events. Type II supernovas are reported to occur with a rate of 0.2 SN units[10)] in spiral galaxies of type *Sb* and 0.57 SN units in *Sc*-galaxies. They have not been observed in *Sa*-, *S0*-, elliptical or irregular galaxies [27]. Therefore we adopt—somewhat arbitrarily—the value

$$x_{\rm ns} \cdot 0.3 \text{ SN units} \tag{26}$$

for the average rate of neutron star (ns) formation. $x_{\rm ns}$ is an unknown factor of order unity.

Then we find from eq.s (25) and (26) a typical time averaged luminosity of flavor f neutrinos from a large sample of stars due to neutron star formation

$$L_{\nu_f}^{\rm ns} = x_{\rm ns} \cdot x_f \cdot 10^{38}\,{\rm sec}^{-1}\,{\rm L}_\odot^{-1}, \tag{27}$$

which compares to the ν_e luminosity of the sun of $1.9\cdot 10^{38}\,{\rm sec}^{-1}\,{\rm L}_\odot^{-1}$. Therefore the time averaged neutrino number-luminosity due to neutron star formation approximately equals the one due to hydrogen burning as one would naively expect.

10) 1 SN unit = 1 Supernova event per $10^{10}\,{\rm L}_\odot$ per 100 years.

If neutrinos of mass m and radiative lifetime τ were produced with fixed energy E in a supernova event in a short pulse of N neutrinos, the expected burst of decay photons at earth would still have a finite width of its energy- and time-distribution. If the supernova occurs at distance $d_{\rm sn}$, the first photons—from neutrino decays near the supernova—would arrive after the time $d_{\rm sn}$, the last ones—from decays near the earth—after the time $d_{\rm sn}/\beta$, where $\beta = \sqrt{1-(m/E)^2}$ is the velocity of the neutrinos. Then the photon pulse duration at the earth would be $\Delta t = T$ where T is defined to be

$$T = d_{\rm sn} \cdot \left(\frac{1}{\beta} - 1\right) \approx \frac{d_{\rm sn}}{2} \cdot \left(\frac{m}{E}\right)^2 . \tag{28}$$

Using electron neutrinos with $m < 46\,\mathrm{eV}$ as an example, taking $E = 10\,\mathrm{MeV}$ and using the above distance to the Virgo cluster for $d_{\rm sn}$, we find $\Delta t \lesssim 2\cdot 10^4\,\mathrm{sec}$. If m is $0.5\,\mathrm{MeV}$ as compatible with the bounds on m_{ν_μ}, one expects $\Delta t \lesssim 2\cdot 10^{12}\,\mathrm{sec} \approx 6\cdot 10^4$ years. In the latter case the pulses from different supernovas would largely overlap and an analysis similar to the one for hydrogen burning stars would be appropriate, using the time averaged neutrino flux (27). As was mentioned before, the limits thus obtainable would not be much better than those from the sun, they would apply, however, to μ-neutrinos, while the solar bounds apply only to electron neutrinos. Assuming that supernovas occur once every two weeks in the Virgo cluster, the photon pulses would be well separated for $m \lesssim 400\,\mathrm{eV}$.

Turning to the case of low masses and well separated pulses we note that the actual photon pulse duration depends on the photon energy and would agree with (28) only for $\omega \approx E$. This is so because only the photons with maximum laboratory energy are the ones which are emitted strictly in the forward direction relative to the neutrino momentum and hence take the shortest path from the supernova to us. All others are emitted at a minute but finite angle. They and their parent neutrinos together trace out two sides of a triangle, the third side of which is the line of sight between the earth and the supernova. To lowest order in $m/E = \gamma^{-1}$, where the photon energies lie in the range $0 \le \omega \le E$, the pulse duration is

$$\Delta t = T \cdot \frac{1 - e^{-2(1-\omega/E)}}{2\left[\frac{\omega}{E} - (\frac{\omega}{E})^2\right]}, \tag{29}$$

where T is the "naive" pulse length (28). The present result expands for $\omega \approx E$ to $T \cdot E/\omega$.

Using the angular distribution (8) for the neutrino decay, the doubly differential photon flux $j_E(\omega,t) = dF/d\omega dt$ is in the limit $m \ll E$

$$j_E(\omega,t) = \frac{N}{4\pi\, d_{\rm sn}^2}\cdot\frac{2}{m\tau}\cdot\left[1+a\left(\frac{2\omega}{E}-1\right)\right]\cdot\frac{\frac{\omega}{E}\cdot e^{-2\omega t/m\tau}}{1-\frac{2t}{T}\cdot\left[\frac{\omega}{E}-\left(\frac{\omega}{E}\right)^2\right]}, \tag{30}$$

where $t = 0$ is defined to be the instant when the first photons arrive, i.e. at the time $d_{\rm sn}$ after the supernova explosion. We note that a typical value for ω is E, for t it is $T = \gamma^{-2}\cdot d_{\rm sn}/2$, therefore the argument of the exponential typically is $-d_{\rm sn}/\gamma\tau$ as one would naively expect. Therefore the exponential can be approximated by unity for $\gamma\tau \gtrsim d_{\rm sn}$. We further note that neutrinos emerging from weak interaction processes would be left-handed, i.e. for $E \gg m$ Dirac neutrinos would be helicity-minus and Dirac anti-neutrinos helicity-plus states. Even for a mixture of these two types of states, as would be expected from a supernova event, the parameter a would take on the definite value -1 as was explained in the discussion following eq. (8).

In a realistic case the neutrinos exhibit a nontrivial energy distribution $f(E)$. If it is normalized, the photon spectrum would then be

$$j_\gamma(\omega,t) = \int_\omega^\infty j_E(\omega,t)\, f(E)\, dE\,. \tag{31}$$

Using the average value $E = 10\,\text{MeV}$ for the neutrino energy, considering the case $a = 0$ and $\omega = E$ and approximating the exponential by unity we find from (30) the doubly differential flux $j_\gamma = 2N/4\pi d_{\rm sn}^2 m\tau$. With the number (25) of neutrinos from a supernova event, using for $d_{\rm sn}$ the distance to the Virgo cluster, taking $m = 10\,\text{eV}$ and $\tau = 10^{11}\,\text{sec}$, as allowed by the solar bounds on electron neutrinos, this is $j_\gamma \approx 10^{-3}\,\text{cm}^{-2}\,\text{sec}^{-1}\,\text{keV}^{-1}$ at $\omega = 10\,\text{MeV}$. It lasts for a time interval $\Delta t \approx 13\,\text{min}$. This pulse would precede an optical supernova event in the Virgo cluster by a day or so because the neutrino pulse precedes the optical outburst by about that time.

We summarize that an x- and γ-ray observation of the Virgo cluster would allow to constrain radiative neutrino lifetimes. Limits on the permanent, continuous radiation would allow to constrain the lifetime of "heavy" neutrinos, $1\,\text{keV} \lesssim m \lesssim 10\,\text{MeV}$, of all flavors. This range could include, at least, the muon-neutrino. The absence of x- and γ-ray bursts in correlation with supernova events would allow to set limits on the lifetime of light neutrinos of all flavors.

This would require a long term observation of the Virgo cluster of typically a few weeks to a few months.

4. Diffuse Neutrino Fluxes from Stellar Sources

4.1 A general limit from diffuse neutrino fluxes

Since the distribution of matter is approximately homogeneous on scales exceeding 50–100 Mpc [28] we are now naturally led to consider the diffuse neutrino flux from the average cosmic distribution of stellar neutrino sources. As was first pointed out by COWSIK [7], the decay of these neutrinos through reaction (1) into radiation would contribute to the diffuse, isotropic x- and γ-radiation. The various measurements of this flux [10] yield a result, an upper limit to which can be conservatively approximated in the range between 100 keV and 10 MeV by the power law representation

$$j_{\text{diffuse}}(\omega) = b \cdot \omega^{-2} = 3 \cdot 10^{-6}\,\text{cm}^{-2}\,\text{sec}^{-1}\,\text{ster}^{-1}\,\text{keV}^{-1} \cdot \left(\frac{1\,\text{MeV}}{\omega}\right)^2. \tag{32}$$

In order to calculate an expected flux which can be compared with this observational result we first consider monochromatic sources of energy E which are homogeneously distributed in space. Hence we consider a neutrino luminosity density $\mathcal{L}_\nu$ which expresses the number of neutrinos which are produced on average per cm^3 and sec in the universe. $\mathcal{L}_\nu$ is assumed to have no intrinsic time dependence and hence varies in an expanding universe like $R^{-3}(t)$, where $R(t)$ is the cosmic scale factor of a homogeneous, isotropic universe [23, 28].

The angular distribution (8) of the decay photons in the neutrino rest frame again determines the energy distribution in the laboratory frame. We adopt the "worst case" with $a = -1$ where the photon spectrum is squeezed furthest toward low frequencies. Then the normalized photon frequency distribution is

$$\frac{dN}{d\omega} = \frac{2}{\omega_{\text{max}}} \cdot \left(1 - \frac{\omega}{\omega_{\text{max}}}\right). \tag{33}$$

Because relativistic neutrinos and photons suffer the same redshift the maximum photon frequency from neutrino-decay of energy E from a source at redshift z is $\omega_{\text{max}} = E/(1+z)$. Using results from ref. [23] we find for the expected, doubly differential photon spectrum $j_\gamma(\omega) = dF_\gamma/d\omega d\Omega$ for the case of large laboratory decay times compared with the age of the universe

$$j_\gamma(\omega) = \frac{\mathcal{L}_\nu}{4\pi} \cdot \frac{m}{\tau} \cdot E^{-2} \int_0^{(\frac{E}{\omega}-1)} dz \cdot 2\left(1 - \frac{(1+z)\omega}{E}\right) \cdot \frac{1}{\dot{R}(z)} \int_{t(z)}^{t(0)} R(t)\,dt\,. \tag{34}$$

The case where the particle ν' is not massless can be simply covered by the substitution $j_\gamma(\omega) \to \frac{1}{r} j_\gamma(\frac{\omega}{r})$ as in our discussion of the solar bounds.

We want to use as few cosmological assumptions as possible. Therefore we consider the universe out to only small redshifts ($z_{\max} = 0.5$) where the expansion can be linearly approximated as [23]

$$\frac{R(t)}{R_0} = \frac{1}{1+z} \approx 1 + H_0 \cdot (t - t_0)\,, \tag{35}$$

where t_0 is the present time, R_0 the present cosmic scale factor and H_0 the present Hubble parameter. We use

$$H_0 = h_0 \cdot 100\,\mathrm{km\,sec^{-1}\,Mpc^{-1}} = (3.08\cdot 10^{17}\,\mathrm{sec}/h_0)^{-1}\,, \tag{36}$$

where observationally [28]

$$0.5 \leq h_0 \leq 1\,. \tag{37}$$

Following WEINBERG [23] we note that, using (35) as an exact relation, errors are introduced into the time evolution of R on the 10–20 % level for $z \leq 0.5$. We further note that galaxies are observed at least out to $z_{\max} = 0.5$ and certainly formed at much larger redshifts. $z_{\max} = 0.5$ means going back in time by $t_{\max} = z_{\max}(1 + z_{\max})^{-1} H_0^{-1} = 10^{17}\,\mathrm{sec}/h_0$ which is slightly less than the age of the sun. We then assume with confidence that stars and galaxies out to such redshifts are similar to the ones at the present time in our region of space.

Then calculating to lowest order in the maximally considered redshift $z_{\max}$ and now assuming a nontrivial (normalized) energy distribution $f(E)$ for the neutrino source spectrum we find for the expected, doubly differential photon flux

$$j_\gamma(\omega) = \frac{\mathcal{L}_\nu}{4\pi} \cdot \frac{m}{\tau} \cdot \frac{1}{2}\left(\frac{z_{\max}}{H_0}\right)^2 \cdot \int_\omega^{E_{\max}} \frac{g(\frac{\omega}{E}, z_{\max}) \cdot f(E)}{E^2}\,dE\,, \tag{38}$$

where

$$g(\tfrac{\omega}{E}, z_{\text{max}}) = 2 \cdot \begin{cases} (1 - \frac{\omega}{E}) & \text{for} \quad 0 \leq \frac{\omega}{E} \leq \frac{1}{1+z_{\text{max}}}, \\ \frac{\omega}{E} (\frac{E}{\omega} - 1)^3 \frac{1}{z_{\text{max}}^2} & \text{for} \quad \frac{1}{1+z_{\text{max}}} \leq \frac{\omega}{E} \leq 1 . \end{cases} \tag{39}$$

is the shape of the photon spectrum for a monochromatic source. It is normalized to first order in z_{max}. The difference between this function and the triangular shape (33) expected from a sample of decaying neutrinos comes from the fact that due to the redshift fewer sources contribute to the high photon energies than to the lower ones.

If we interpret, to lowest order in z_{max}, $r_{\text{max}} = z_{\text{max}}/H_0$ as the radius of the farthest spherical shell of the universe from which radiation is observed and $t_{\text{max}} = r_{\text{max}}$ as the time of travel of a relativistic particle from there to us, and considering a monochromatic neutrino source, then the doubly integrated photon flux simply is

$$F_\gamma = \frac{t_{\text{max}}/2}{\tau_{\text{lab}}} \cdot \mathcal{L}_\nu \, r_{\text{max}}, \qquad \tau_{\text{lab}} = \frac{E}{m} \tau \, , \tag{40}$$

which compares to eq. (5) of ref. [7].[11]

From the requirement $j_\gamma(\omega) < j_{\text{diffuse}}(\omega)$ and from the power law representation (32) for j_{diffuse} we find the general bound

$$\frac{\tau}{m} > \frac{8}{27} \cdot \frac{\mathcal{L}_\nu}{4\pi} \cdot \frac{1}{b} \cdot \frac{1}{2} \Big(\frac{z_{\text{max}}}{H_0} \Big)^2 \cdot \max_{0 \leq \omega \leq E_{\text{max}}} \frac{27}{8} \int_{\omega}^{E_{\text{max}}} \Big(\frac{\omega}{E} \Big)^2 g(\tfrac{\omega}{E}, z_{\text{max}}) \, f(E) \, dE \, . \tag{41}$$

For a monochromatic neutrino source, $f(E) = \delta(E_0 - E)$, and for $z_{\text{max}} = 0.5$ the last factor is unity, independently of E_0. For other energy distributions it is slightly smaller.[12] For $z_{\text{max}} = 0.5$ the first part of this formula is numerically $2.80 \cdot 10^{39}\,\text{sec/eV} \cdot \mathcal{L}_\nu\,\text{cm}^3\,\text{sec} \cdot h_0^{-2}$.

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[11] COWSIK misses, however, the factor 1/2 which is due to the fact that *on average* neutrinos have only the time $t_{\text{max}}/2$ available for their decay.

[12] When $f(E)$ is the allowed β-decay spectrum of neutrinos from the pp-reaction with an average energy of 265 keV it is 0.62. When f is a similar spectrum for the neutrinos expected from ^{8}B decay with an average energy of 7.1 MeV which is important in the *PPIII*-chain, it is 0.58. When f exhibits several peaks—a case which we will consider in the first application of this formula—this factor will be accordingly smaller.

4.2 Hydrogen burning stars

In order to derive definite limits we must consider definite neutrino sources. In the case of hydrogen burning stars we translate the bolometric luminosity density of the universe into a neutrino luminosity in exactly the same fashion as in our discussion of the neutrino spectrum emitted by the galactic center. We base our treatment on the measured luminosity density [29] from galaxies of

$$\mathcal{L}_{\rm bol} = 1.8{\cdot}10^{-26}\,{\rm MeV\,sec^{-1}\,cm^{-3}} \cdot h_0\,. \tag{42}$$

We attribute 75 % of this energy production to hydrogen burning.

Varying the relative rates of the four relevant reaction chains we find a result almost identical to the one in Sect. 3.1 for the relative importance of the four reaction chains. This result means that 2 neutrinos are emitted for each 22.1 MeV of bolometric energy release of stars due to hydrogen burning. Then we find for the average cosmic neutrino production due to hydrogen burning the value

$$\mathcal{L}_{\nu_e}^{\rm hb} = 1.2{\cdot}10^{-27}\,{\rm sec^{-1}\,cm^{-3}} \cdot h_0\,. \tag{43}$$

The last factor in our general formula (41) turns out to be 0.30. The smallness of this number compared with unity is due to the fact that the neutrino spectrum in the present case exhibits two peaks of about equal magnitude.

Then we obtain the limit

$$\tau_{\nu_e}/m_{\nu_e} > 1.0{\cdot}10^{12}\,{\rm sec/eV} \cdot h_0^{-1}\,. \tag{44}$$

This result can be translated into limits for the subdominantly coupled neutrino admixtures in much the same way as has been described in the context of our solar bounds.

4.3 Supernovas and the formation of neutron stars

In their late stages of evolution stars tend to form degenerate ("white") dwarfs or—if they are initially heavy enough—neutron stars. Turning first to the latter case we find from eq.'s (27) and (42) for the average production rate of neutrinos of flavor f from supernova events

$$\mathcal{L}_{\nu_f}^{\rm sn} = x_{\rm ns} \cdot x_f \cdot 7.5{\cdot}10^{-28}\,{\rm sec^{-1}\,cm^{-3}} \cdot h_0\,. \tag{45}$$

We estimate the last factor in our formula (41) to be 0.5 and hence find the neutrino radiative lifetime bound

$$\tau_{\nu_f}/m_{\nu_f} > x_{\rm ns} \cdot x_f \cdot 1.0{\cdot}10^{12}\,{\rm sec/eV} \cdot h_0^{-1}\,, \tag{46}$$

identical to the one from hydrogen burning stars, but applicable, at least, also to μ-neutrinos which are known to be lighter than 0.5 MeV.

COWSIK in his original treatment [7] finds instead the bound $3{\cdot}10^{16}\,{\rm sec/eV}$. His assumptions lead to a total neutrino production from supernovas of $\mathcal{L}_\nu^{\rm sn} = 1.3{\cdot}10^{-26}\,{\rm sec^{-1}\,cm^{-3}}$. Taking from eq. (32) $3{\cdot}10^{-8}\,{\rm cm^{-2}\,sec^{-1}\,ster^{-1}\,keV^{-1}}$ as a flux limit at 10 MeV, assuming a "box-spectrum" for the photons and using formula (40) for the total photon flux with COWSIK's values of $r_{\rm max} = 3{\cdot}10^{28}\,{\rm cm}$ and $t_{\rm max} = 10^{18}\,{\rm sec}$ we find $\tau_\nu/m_\nu > 5{\cdot}10^{15}\,{\rm sec/eV}$. We believe that this is the result that would be consistent with his assumptions.

COWSIK's choice for the cutoff-in-the-past, however, appears to be inconsistent with his neglect of cosmological redshifts and with established values for the age of the universe which are about a factor of 2–3 below $10^{18}\,{\rm sec}$. We believe that going further into the past than our value $t_{\rm max} = 10^{17}\,{\rm sec}$ requires a more detailed discussion of redshifts and cosmological models than has been presented either by COWSIK or in the present note. The final result is proportional to the square of $t_{\rm max}$, therefore COWSIK's result should be further reduced by two orders of magnitude.

COWSIK's estimate of the average neutrino production is more than an order of magnitude above our estimate. Although his value as well as ours are rather gross estimates and can therefore easily be disputed, we feel that we have tried to find a conservative but fair value which is in agreement with published data on supernova events [27]. We believe, then, that COWSIK's result should be further reduced by an order of magnitude if one aims at a conservative bound as we have attempted to do.

Then the residual discrepancy of about a factor of 6 between COWSIK's and our result can be attributed to the fact that we have treated the relevant photon spectra in more detail than he has done.

4.4 Formation of degenerate dwarfs

Stars with masses of up to a few solar masses will in their late stages of evolution collapse to form degenerate dwarfs of mass not larger than the Chandrasekhar limit of about 1.4 $M_\odot$. The remainder of the mass is ejected and forms a planetary nebula. Following COWSIK [7] we assume that most of the gravitational binding energy of the white dwarf of about $GM_{\rm wd}^2/R_{\rm wd}$ is radiated away in the form of 100 keV neutrinos. Considering $M_{\rm wd} = M_\odot = 2{\cdot}10^{33}\,$g and $R_{\rm wd} = 0.006R_\odot = 4{\cdot}10^8\,$cm [30] this corresponds to about $4{\cdot}10^{56}\,$MeV and thus to the production of about $4{\cdot}10^{57}$ neutrinos. This number would be considerably smaller for less massive white dwarfs because their mass is smaller and their radius is *larger*. For comparison we note that 1 $M_\odot$ of pure hydrogen, when totally burnt into helium, releases about $8{\cdot}10^{57}\,$MeV of nuclear binding energy and $6{\cdot}10^{56}$ neutrinos. This means that a star in its evolution emits at the most about an order of magnitude more neutrinos during the formation of a white dwarf than during its hydrogen burning stages.

Neutrino radiative lifetime limits would then be more restrictive by about an order of magnitude than the limits obtained from hydrogen burning stars.

COWSIK's corresponding limit [7] is $10^{17}\,$sec/eV which we find to be consistent with his assumptions. However, he uses $\mathcal{L}_\nu^{\rm wd} = 3{\cdot}10^{-25}\,{\rm sec}^{-1}\,{\rm cm}^{-3}$, about two orders of magnitude above our rate from hydrogen burning stars. We believe that this number should be reduced by an order of magnitude according to our above estimate. Furthermore a similar criticism applies to this result as in the case of neutron star formation. Therefore this result should be reduced by 3–4 orders of magnitude.

5. Summary

We have derived in a detailed analysis what can be learned from astrophysical neutrino and photon fluxes about the radiative decay of neutrinos. x- and γ-ray measurements of the sun and of the galactic center constrain the lifetime-over-mass-ratio of the electron neutrino to $\tau_{\nu_e}/m_{\nu_e} > 10^9 - 10^{10}\,$sec/eV. Measurements of the diffuse cosmic background flux of x- and γ-rays yield a bound $\tau_\nu/m_\nu > 10^{12}\,$sec/eV, applicable to any neutrino species lighter than about 10 MeV. We have speculated on potential bounds obtainable from observations of supernovas in the Virgo cluster. We believe that at least the solar bounds are practically as reliable as the available laboratory results.

These numbers are far below any theoretical predictions of the standard model for the case of neutrino mixing. Only bounds which are based on cosmological neutrino sources could possibly come close to such predicted values. The following appendix gives a brief summary of such limits which have been discussed in the literature. The significance of cosmological results in comparison with laboratory measurements, however, is not generally agreed upon. The intent of the present discussion was, therefore, to derive astrophysical bounds which would parallel laboratory results as closely as possible—both, in method and significance.

Appendix on Neutrino Radiative Lifetime Constraints from Astrophysical and Cosmological Arguments

We briefly review the limits on the lifetime of massive neutrinos that have been obtained from astrophysical and cosmological arguments by various authors. All of these limits are based on *thermal* neutrino sources. Therefore charged current neutrino-producing reactions are not particularly important, hence the flavor-identity of the neutrinos is rather irrelevant. The following arguments are therefore applicable to any species of weakly interacting neutral fermions.

Neutrino production from the heat bath in the interior of *white dwarfs* has been used by COWSIK [7] and GOLDMANN & STEPHENSON [31] to calculate the expected contribution to the cosmic diffuse x- and γ-ray background from neutrino decay. From observational bounds on these backgrounds they derive a limit $\tau_\nu/m_\nu \gtrsim 10^{17}\,\mathrm{sec/eV}$ for $m_\nu \lesssim 100\,\mathrm{keV}$.[13]

Supernova-Explosions were used in a similar way by COWSIK [7, 32]. He obtained $\tau_\nu/m_\nu \gtrsim 3\cdot 10^{16}\,\mathrm{sec/eV}$ for $m_\nu \lesssim 10\,\mathrm{MeV}$.[13] FALK & SCHRAMM [33] consider, in addition, details of supernova energetics for the case where the neutrinos decay within the presupernova star. They exclude a region of the m_ν-τ_ν-plane in a range $10^{-3}\,\mathrm{sec} \lesssim \tau_\nu \lesssim 10^3\,\mathrm{sec}$, $10\,\mathrm{keV} \lesssim m_\nu \lesssim 10\,\mathrm{MeV}$.

All of the following cosmological bounds are based on the flux of neutrinos which is generally believed to have been produced in the *Big Bang of the Universe*. They thus depend on the assumption that standard cosmology correctly describes the history of our universe.

If these primordial neutrinos reside now in the dark halos of galaxies or of galactic clusters they would probably have masses of a few tens of eV. Their

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[13] Consider, however, our criticism of these numbers in the main body of the present paper.

radiative decays would produce a photon flux in the ultraviolet range. From the observed background fluxes or from fluxes of discrete objects such as the Coma- or Virgo-cluster, bounds in the typical range $\tau_\nu \gtrsim 10^{22} - 10^{25}$sec have been obtained and/or discussed by HENRY & FELDMAN [34], by HOLBERG & BARBER [35], by KIMBLE, BOWYER & JAKOBSEN [36], by DE RUJULA & GLASHOW [37], by SHIPMAN & COWSIK [38], by STECKER [39] and by WELLER [40].

Neutrinos produced in the Big Bang must decay so fast that the decay photons can be thermalized before radiation decouples from matter, otherwise the cosmic microwave background would be distorted. They could also decay so slowly that they do not provide a photon signal in the uv-, x- and γ-range that would violate observed bounds on the diffuse background radiation spectrum. Using such arguments, large regions in the τ_ν-m_ν-plane have been excluded by COWSIK [7], by DICUS, KOLB & TEPLITZ [41], by GUNN, LEE, LERCHE, SCHRAMM & STEIGMAN [42], by MCKELLAR & PAKVASA [43] and by SALATI [44]. The argument was extended by SILK & STEBBINS [45] to the case where thermalization is incomplete before the universe becomes transparent. They consider lifetimes in the range $10^8\,\text{sec} \lesssim \tau_\nu \lesssim 10^{12}\,\text{sec}$.

If the decay photons are energetic enough they would have dissociated hydrogen in conflict with the observation of neutral hydrogen clouds in the Milky Way and other galaxies. MELOTT & SCIAMA [46] and REPHAELI & SZALAY thus find $\tau_\nu \gtrsim 10^{24}\,\text{sec}$ for $30\,\text{eV} \lesssim m_\nu \lesssim 150\,\text{eV}$.

Re-ionization of the intergalactic medium after the cosmic era of recombination by radiatively decaying "-inos" has been discussed by SALATI [47].

The effect of unstable neutrinos on primordial nucleosynthesis has been considered by DICUS, KOLB, TEPLITZ & WAGONER [48], by LINDLEY [49], by MIYAMA & SATO [50] and by SATO & KOBAYASHI [51]. From data on what is believed to be the primordial abundances of light elements and the present baryon-to-photon-ratio they exclude large areas in the mass-lifetime-plane for the neutrinos.

A summary of many of the preceding bounds has been given, e.g. by BARROW [28], by DOLGOV & ZELDOVICH [17] and by TURNER [52]. These authors also review the bounds on neutrino masses that can be obtained from limits on the present energy density of the universe. Such bounds exclude certain areas in the neutrino mass-lifetime-plane, the emphasis being on mass-limits for sufficiently long lived neutrinos and not on the hypothetical radiative decay mode.

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