NEUTRINOS

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In addition to the ν_e, ν_μ, and ν_τ sections, the *Review of Particle Physics* includes sections on "Number of Light Neutrino Types," "Heavy Lepton Searches," and "Searches for Massive Neutrinos and Lepton Mixing."

Neutrino experiments are notoriously difficult, owing to the basic property that neutrinos are neutral, weakly interacting particles. Over the years, many experimental claims pertaining to neutrino properties have been refuted by subsequent data. The *Review of Particle Physics* is an archival compendium which includes references to older papers, even experimental claims which have now been ruled out. It will be clear from the various listings which experiments have later been refuted.

In view of the continuing unsettled nature of data pertaining to various neutrino properties, it is perhaps well to record some of the definite accomplishments in the history of the subject. Neutrinos were first proposed in 1930 by Pauli, to explain the observed continuous electron energy distribution in nuclear beta decay [1]. Tentative evidence for the observation of the electron (anti)neutrino was reported in 1953, and definite evidence in 1956, by Cowan, Reines, and coworkers, using the reaction $\overline{\nu}_e p \to e^+ n$ with $\overline{\nu}_e$'s from reactor fluxes [2]. The separate identity of ν_e and ν_μ was demonstrated experimentally in 1962 by Lederman, Schwartz, Steinberger, and coworkers in a Brookhaven experiment [3]. Neutrinos from the sun were first observed by R. Davis and coworkers via the reaction $\nu_e\ {}^{37}\mathrm{Cl} \to e^-\ {}^{37}\mathrm{Ar}$, using an underground radiochemical experiment (which began operation in the late 1960's) in the Homestake Gold Mine [4]. Although we tabulate here only experiments with results pertaining to neutrino properties, it should be recalled that neutrino reactions played a crucial role in confirming the now-Standard Model when in 1973 neutral weak currents were first observed by the CERN Gargamelle bubble chamber experiment via the reactions $\nu_\mu e \to \nu_\mu e$ and $\nu_\mu(\overline{\nu}_\mu) + N \to \nu_\mu(\overline{\nu}_\mu)$ *hadrons* [6]. Neutrino reactions have also provided an important input to the measurement of the weak mixing angle θ_W. The discovery of the τ lepton by Perl *et al.* at

SPEAR in 1975 and the study of its decay implied the existence of the third neutrino, ν_τ [7]. The precise measurement of the width of the Z at LEP and SLC has shown that there are only three species of neutrinos (in the usual electroweak doublets) with masses $< m_Z/2$ [8].

The theoretical perspective concerning neutrino masses has changed considerably over the past 20 years. Before that time, a standard view was that there was no theoretical reason for neutrinos to have masses, which was in accord with the striking fact that the upper limits on their masses were much smaller than those of the associated charged leptons. It was also noted that experimental data were consistent with the "laws" of lepton family number and total lepton number conservation. (Some early discussions of of neutrino oscillations and lepton mixing are given in Refs. 9 and 10). In the literature through the 1970's one often finds statements asserting that in the standard SU(2) $\times$ U(1) electroweak theory (without electroweak-singlet neutrinos) the known neutrinos (in electroweak-doublets) are massless. This is true if one pretends that the Standard Model is applicable to arbitrarily high energies and requires the exact absence of any nonrenormalizable, higher-dimension operators in this theory. However, a more modern view is that the Standard Model is an effective field theory, which is a good description of nature only up to some energy scale where new physics occurs. Clearly a strict upper bound on this scale is given by the Planck mass, $\overline{M}_{Pl} \equiv \sqrt{\hbar c/(8\pi G_N)} = 2.4 \times 10^{18}$ GeV, since quantum gravity is not included in the Standard Model. However, there are strong arguments that new physics beyond the Standard Model actually occurs at a much lower scale, of order a TeV. This new physics may be able to be included in a generalization of the Standard Model which remains perturbative, as in supersymmetric extensions, or may be nonperturbative, as in dynamical electroweak symmetry breaking schemes. It has been appreciated that renormalizability and, in particular, the great success of the Standard Model with its exclusion of any higher-dimension nonrenormalizable operators, may well be due only to the fact that the electroweak scale v_{EW} is considerably smaller than the scale of new physics. A summary of this modern view is given, *e.g.*, in Ref. 11.

Once one includes higher-dimension operators in the Lagrangian, nonzero neutrino masses can easily occur. The sizes of these masses reflect the scale(s) of the new physics. For example, given only the known left-handed neutrino fields of the Standard Model, there would be a gauge-invariant dimension-5 operator

$$\mathcal{O} = \frac{1}{M_X} \sum_{a,b} h_{a,b}(\epsilon_{ik}\epsilon_{jm} + \epsilon_{im}\epsilon_{jk}) \left[\mathcal{L}^{Ti}_{La} C \mathcal{L}^{j}_{Lb}\right] \phi^k \phi^m + h.c. \quad (1)$$

where $\mathcal{L}_{La} = (\nu_{\ell_a}, \ell_a)^T_L$ is the left-handed, $I = 1/2$, $Y = -1$ lepton doublet with generation index a ($a = 1$, 2, or 3), where $\ell_a = e, \mu, \tau$, for $a = 1, 2, 3$, and M_X denotes a generic mass scale characterizing the origin of this term. This operator involves a symmetric, $I = 1$ combination of the two $I = 1/2$ lepton doublets, contracted with an $I = 1$ combination of the two Higgs doublets. The term arising from the vacuum expectation values (vev's) of the Higgs doublets yields a (left-handed) Majorana neutrino mass term (symmetric in generation indices). If $M_X >> v_{EW}$, where $v_{EW} = 2^{-1/4}G_F^{-1/2} = 246$ GeV is the electroweak symmetry breaking scale, this would explain the smallness of the resultant neutrino masses.

Because of the hierarchy problem plaguing the Higgs sector of the Standard Model, many physicists have concluded that either this sector is stabilized against large radiative corrections by supersymmetry, or the electroweak symmetry breaking originates not from the vacuum expectation value of a pointlike Higgs field, but instead dynamically, from the condensation of bilinear products of fermion fields without any fundamental Higgs field. In supersymmetric extensions of the Standard Model, one again finds dimension-5 operators analogous to Eq. (1). In approaches based on dynamical electroweak symmetry breaking, one can also get neutrino mass terms arising from higher-dimension multifermion operators.

In contrast to the higher-dimension operator (1), which involves only the known neutrinos, together with the hypothetical Higgs field of the Standard Model (or its supersymmetric extensions), another mechanism makes use only of renormalizable, dimension-4 operators, but requires the existence of electroweak-singlet neutrino fields. As will be discussed below, this mechanism produces light neutrino masses of order $m_\nu \sim M_D^2/M_R$, where M_D denotes a generic Dirac neutrino mass and M_R denotes a generic electroweak-singlet Majorana neutrino mass. Since the mass scale of the electroweak singlet neutrino mass term is naturally $>> v_{EW}$, this again yields, albeit for a different reason, very small m_ν [12].

In turn, a natural concomitant of (nondegenerate) neutrino masses is lepton mixing, which is thus also a general expectation. The lepton-mixing angles are functions of ratios of elements of neutrino-matrix elements and of charged lepton mass matrix elements, and even though left-handed neutrino masses are small, some of these ratios could, in principle, be $\mathcal{O}(1)$, which raises the issue of why such effects have not been seen. This question was answered as follows: a set of conditions for natural suppression of observable lepton flavor violation were formulated, and it was shown that the Standard Model (generalized to include nonzero m_ν) satisfies these [13]. This explains why the "law" of lepton family number conservation is obeyed to such high accuracy.

After these theoretical points, let us return to a description of the quantities upon which various experiments put limits. As an aid to understanding the limits on neutrino masses and lepton mixing, we recall that, in contrast to other particles in this *Review*, the neutrinos ν_e, ν_μ, and ν_τ are defined as weak eigenstates (the weak $I_3 = 1/2$ components of the $\mathrm{SU}(2)_L$ lepton doublets) which couple with unit strength to e, μ, and τ, respectively. These neutrino weak eigenstates are not, in general, states of definite mass. If one assumes that neutrinos are massless, and hence degenerate, then it is possible to define the weak eigenstates to be simultaneously mass eigenstates. However, in the general case of possibly massive (nondegenerate) neutrinos, the weak eigenstates have no well-defined masses, but instead are linear combinations of mass eigenstates. Let us denote the charged leptons as the set $\{\ell_a\}$, $a = 1, \ldots, n$, where $n \geq 3$, with $\ell_1 = e$, $\ell_2 = \mu$, and $\ell_3 = \tau$. From the LEP measurement of the Z width (see section on "Number of Light Neutrinos"), one knows that there are only three neutrinos which couple to the Z in the usual way and have masses $m_\nu < m_Z/2$. Of course, this measurement does not preclude the existence of electroweak-singlet neutrinos. The left-handed components of the weak eigenstates of the neutrinos, $(\nu_{\ell_a})_L$ can be expressed in terms of mass eigenstates by the transformation

$$(\nu_{\ell_a})_L = \sum_j U_{aj}(\nu_j)_L \tag{2}$$

where the $\{\nu_j\}$ denote these mass eigenstates. The mass eigenstates are, in general, linear combinations of the known $I = 1/2$, $I_3 = 1/2$ neutrinos from electroweak doublets, and, in addition, possible electroweak-singlet neutrinos (sometimes called "sterile" neutrinos). The ordering of the mass eigenbasis can be defined so that U is as nearly diagonal as possible, *i.e.* (with no sum on j) $|U_{jj}| \geq |U_{jk}|$, $k \neq j$. Of course, this does not imply that $m(\nu_j) > m(\nu_k)$ for $j > k$.

Thus, as was noted in Ref. 14, decays such as ${}^3\mathrm{H} \to {}^3\mathrm{He}\, e^- \, \overline{\nu}_e$ and $\pi^+ \to \mu^+\nu_\mu$, which have been used to set the best bounds on the respective neutrino masses, really consist of sums of the separate decay modes ${}^3\mathrm{H} \to {}^3\mathrm{He}\, e^- \, \overline{\nu}_j$ and $\pi^+ \to \mu^+\nu_k$, where the ν_j and ν_k are mass eigenstates, and the indices j and k range over all of the values allowed by phase space in these respective decays. The coupling strengths for the jth mode in ${}^3\mathrm{H}\beta$ decay and the kth mode in $\pi^+_{\mu 2}$ decay are given, respectively, by $|U_{1j}|^2$ and $|U_{2k}|^2$. In general, these modes are incoherent, although in the limit in which the ν_j all become degenerate they would become coherent. There are, in addition certain kinematic factors depending on the m_{ν_j} which enter in determining the branching ratio for a given decay mode. Assuming that the off-diagonal elements of the lepton mixing matrix U are small relative to the diagonal elements, the dominantly coupled decays are the ones with coupling strength $|U_{aj}|^2$, $a = j$, *i.e.*, ${}^3\mathrm{H} \to {}^3\mathrm{He}\, e^- \, \overline{\nu}_1$ and $\pi^+ \to \mu^+\nu_2$.

Hence, it follows that the neutrino mass limits quoted in the literature for "m_{ν_e}," "m_{ν_μ}," and "m_{ν_τ}" should really be interpreted as limits on the corresponding mass eigenstates [14,15]. Specifically, a bound on "m_{ν_e}" from a study of tritium β decay, for example, really constitutes a weighted limit on each of the mass eigenstates ν_j in the weak eigenstate ν_e which are kinematically allowed to occur in tritium decay and which are coupled with strength $|U_{1j}|^2$ sufficiently large to make a significant contribution to the observed spectrum. It is thus certainly a limit on ν_1, since this is, by the definition, of the order of the mass eigenbasis, the dominantly coupled neutrino. If lepton mixing is hierarchical, as quark mixing is known to be, *i.e.*, if $|U_{jj}|^2 \gg |U_{jk}|^2$, $j \neq k$, then ν_1 is the only mass eigenstate significantly constrained by a bound on "m_{ν_e}." Furthermore, strictly speaking, a neutrino mass limit cannot be stated in isolation; it always contains some implicit dependence on the relevant lepton-mixing angles. This dependence is fortunately relatively unimportant for the dominantly coupled decay modes, *i.e.*, $e\,\overline{\nu}_1$, $\mu\,\overline{\nu}_2$, and $\tau\,\overline{\nu}_3$ and hence the mass limits on "m_{ν_e}," "m_{ν_μ}," and "m_{ν_τ}" can be reinterpreted as being limits on m_{ν_j}, $j = 1,\ 2$, and 3, respectively.

There are three general types of (Lorentz-invariant) neutrino mass terms: Dirac masses of the form $m_D\,\overline{\nu}_L\,\chi_R + h.c.$, left-handed Majorana masses of the form $m_L\,\overline{\nu}_L\,\nu^c_R + h.c. = m^*_L\,{\nu_L}^T\,C\nu_L + h.c.$ and right-handed Majorana masses of the form $m_R\,\overline{\chi}^c_L\,\chi_R + h.c. = m_R\,\chi_R^T C\,\chi_R + h.c.$, where C is the Dirac charge conjugation matrix. Clearly, Dirac and right-handed Majorana mass terms require the existence of electroweak-singlet neutrinos. Our notation χ_R follows the usual practice of calling these "right-handed neutrino singlets", although since they are singlets, it is a convention whether one writes them as χ_R or $\chi'_L = (\chi_R)^c$. It is not known whether such electroweak-singlet neutrinos actually exist. Dirac mass terms conserve total lepton number $L_{\rm tot}$, while Majorana mass terms violate $L_{\rm tot}$. In the standard electroweak theory, extended to include massive neutrinos, (i) a Dirac mass term transforms as a weak $I = 1/2$ operator, and is coupled to the $I = 1/2$ Higgs to make an SU(2) $\times$ U(1) singlet operator; (ii) a Majorana mass term involving the $I = 1/2$ left-handed neutrinos transforms as $I = 1$ and must be coupled to an operator with $I = 1$ (and $Y = 2$) to make a gauge-invariant singlet; (iii) a Majorana mass term involving the SU(2) $\times$ U(1) singlet neutral leptons, conventionally considered to be right-handed, is a singlet; it could be present as a bare mass term or couple to some other singlet operator. Note that in the minimal supersymmetric Standard Model (MSSM), which has two Higgs doublets, of hypercharge $Y = 1$ and $Y = -1$, the Dirac neutrino mass term arises from the cubic chiral superfield terms $\epsilon_{ij} \sum_{a,b} \hat{L}^i_a \hat{\chi}^c_b \hat{H}^j_u$ (all chiral superfields are taken as left-handed), where $\hat{H}_u$ is the same Higgs that gives mass to the $Q = 2/3$ quarks. The Dirac neutrino mass terms are thus proportional to $\sin\beta$, where $\tan\beta = v_u/v_d$ is the ratio of the vacuum expectation values of the two Higgs in the MSSM.

In general, in the Standard Model, in addition to the three known left-handed $I = 1/2$ lepton doublets, there could be some number n_s of electroweak-singlet neutrinos. In a compact notation, one can then denote ν_L as the 3-component vector of left-handed $I = 1/2$ neutrinos and χ_R to be the n_s-dimensional vector of electroweak-singlet singlets, taken to be right-handed. The general neutrino mass term in the Lagrangian is then given by

$$-\mathcal{L}_m = \frac{1}{2}(\overline{\nu}_L, \overline{\chi}^c_L) \begin{pmatrix} M_L & M_D \\ M_D^T & M_R \end{pmatrix} \begin{pmatrix} \nu^c_R \\ \chi_R \end{pmatrix} + h.c. \qquad (3)$$

where M_L is the 3×3 left-handed Majorana-mass matrix, M_R is a $n_s \times n_s$ right-handed Majorana-mass matrix, and M_D is the 3-row by n_s-column Dirac-mass matrix. In general, all of these mass matrices are complex. The anticommutativity of fermion fields and the property that $C\gamma_\mu C^{-1} = -{\gamma_\mu}^T$ together imply that the Majorana mass matrices are symmetric: $M_L = {M_L}^T, M_R = {M_R}^T$. The diagonalization of the full $(3 + n_s) \times (3 + n_s)$ mass matrix in Eq. (2) yields $3 + n_s$ mass eigenstates, which are, in general, of Majorana type. Since Majorana mass terms violate total lepton number, one sees from a general viewpoint that one does not expect conservation of total lepton number. In particular, the dimension-5 operators discussed above give rise to left-handed Majorana neutrino mass terms and violate total lepton number. Dirac-neutrinos can be constructed from two Majorana-neutrino mass eigenstates whose masses are equal in magnitude [16]. For this reason, Dirac neutrino masses may be considered to be a special (degenerate) case of Majorana neutrino masses, and the latter may be regarded as the generic case. From the transformation which diagonalizes the neutrino mass matrix, together with the transformation which diagonalizes the charged lepton mass matrix (where, of course, only Dirac masses are allowed by electric-charge conservation), one constructs the lepton-mixing matrix U. In general, since U is not the identity, neutrino masses naturally give rise to lepton family number violation.

In supersymmetric extensions of the Standard Model, the neutrinos could, *a priori*, mix with the neutralinos (higgsinos and neutral gauginos). However, the usual R parity which is invoked to forbid unacceptably rapid proton decay also prevents such mixing between neutrinos and neutralinos.

In addition to mass and lifetime limits, this *Review* includes limits on various other possible properties, including electric charge, the CPT-violating difference $m_{\nu_1} - m_{\overline{\nu}_1}$, and a magnetic dipole moment. These are of interest because a massless purely chiral Dirac neutrino cannot have a magnetic (or electric) dipole moment. In the standard electroweak theory, extended to allow for Dirac neutrino masses, the neutrino magnetic dipole moment is nonzero and given [13,17], as

$$\mu_{\nu_j} = \frac{3eG_F m_{\nu_j}}{8\pi^2\sqrt{2}} = 3.2 \times 10^{-19}(m_{\nu_j}/1 \text{ eV})\mu_B \tag{4}$$

where G_F is the Fermi constant and $\mu_B = e/2m_e$ is the Bohr magneton. The neutrino electric dipole moment violates both time-reversal invariance and parity; although it is nonzero in general, it is quite small (see, *e.g.* Ref. 18). Again, however, we note that Dirac neutrinos should be regarded as a special case; the generic case is Majorana neutrinos. Because of the properties

$$C\Gamma C^{-1} = -\Gamma \ , \qquad \Gamma = \sigma_{\alpha\beta} \ , \sigma_{\alpha\beta}\gamma_5 \tag{5}$$

it follows that the operator products which define the magnetic and electric dipole moments are of the respective forms

$$(\overline{\nu}_i \sigma_{\alpha\beta} \nu_j - \overline{\nu}_j \sigma_{\alpha\beta} \nu_i) F^{\alpha\beta} \tag{6}$$

and

$$(\overline{\nu}_i \sigma_{\alpha\beta} \gamma_5 \nu_j - \overline{\nu}_j \sigma_{\alpha\beta} \gamma_5 \nu_i) F^{\alpha\beta} \tag{7}$$

(where $F^{\alpha\beta}$ is the electromagnetic field strength tensor). Hence, if ν_i is a Majorana neutrino (mass eigenstate), its magnetic and electric dipole elements vanish identically. Although only the diagonal magnetic and electric dipole moments are static properties of a given neutrino mass eigenstate, transition magnetic and electric dipole moments may exist, in general, for both Dirac and Majorana neutrinos.

Occasionally, one also finds references to the "neutrino charge radius" in the literature. This is defined via the Taylor series expansion of the generalized vector Dirac form factor multiplying γ_μ in the electromagnetic current matrix element: $F_1^V(q^2) = F_1^V(0) + q^2\, dF_1^V/dq^2|_{q^2=0} + \mathcal{O}[(q^2)^2]$, where q denotes the 4-momentum of the photon [see, *e.g.* Ref. 13 Eq. (2.20)]. The electric charge is $Q = F^V(0) = 0$ for a neutrino, and the charge radius is given by $\langle r^2 \rangle = (1/6)(F_1^V)'(0)$. However, since this is multiplied by q^2 in the Taylor series expansion, it never occurs for a real photon, where $q^2 = 0$, and hence is not an S-matrix element, *i.e.*, not a physical quantity. In a gauge theory, this is manifested in the fact that the charge radius is gauge-dependent.

If one considers the possibility of nonzero masses for neutrinos, for consistency one must then also consider the leptonic mixing which would in general occur concomitantly. Accordingly, this *Review* devotes a section to correlated bounds on neutrino masses and lepton mixing angles. These can be divided into two types. First, there are those due to decays involving neutrinos in the final state, which must be recognized to have the possible multimode structure pointed out above. In the two most sensitive cases suggested as tests for neutrino masses and mixing, one obtains a limit on m_{ν_j} and $|U_{aj}|^2$ individually for each j. The peak-search test proposed in Ref. 14 was applied to existing data in that paper and a subsequent one [15]; it was applied in new experiments on 2-body leptonic decays of K^+ and π^+ by several groups at SIN (PSI), KEK, and TRIUMF. The results are catalogued in corresponding subsections on limits on $|U_{1j}|^2$ and $|U_{2j}|^2$. The kink-search test was also applied by a number of groups. The experimental situation, which was controversial for many years, has recently been clarified (see below).

Second, there are those due to processes involving the propagation and subsequent interaction of neutrinos. The latter are often called neutrino-oscillation limits, although this term is strictly correct only if the differences in neutrino masses are sufficiently small relative to their momenta that the propagation is effectively coherent in a quantum mechanical sense; otherwise, the individual ν_j from a given decay such as $\pi_{\mu 2}$ or $K_{\mu 2}$ propagate in a measurably incoherent manner, and there is no oscillation. Experimentalists usually present their results in terms of a simplifying model in which mixing is assumed to occur only between two neutrino species. The relevant transformation equation becomes

$$\begin{pmatrix} \nu_{\ell_a} \\ \nu_{\ell_b} \end{pmatrix} = \begin{pmatrix} \cos\theta & \sin\theta \\ -\sin\theta & \cos\theta \end{pmatrix} \begin{pmatrix} \nu_i \\ \nu_j \end{pmatrix} \tag{8}$$

where ν_{ℓ_a} are the neutrino weak eigenstates, with $\nu_{\ell_1} = \nu_e$, *etc.*, and ν_i are neutrino mass eigenstates. A given decay, such as $\pi^+ \to \mu^+\nu_\mu$, produces, at time $t = 0$, the neutrino mass eigenstates which are contained in the weak eigenstate ν_μ and are kinematically allowed to occur as decay products. Each mass eigenstate picks up a phase as it propagates, so that at time t, the jth eigenstate, as a quantum-mechanical state, has acquired a phase $\exp(-iE_jt)$, where E_j denotes its energy. Strictly speaking, E_j must be considered to be complex, since a massive neutrino will, in general, decay. Indeed such neutrino decays have been searched for in various experiments (see Listings). In the present discussion, we shall neglect this; for the ranges of neutrino masses of relevance to terrestrial neutrino oscillation experiments, such decays should have a negligible effect on the observed oscillations. (Indeed, the observation of neutrinos at about the expected rate from the 1987 supernova places significant lower bounds on neutrino lifetimes.) According to basic quantum mechanics, one cannot measure the energies E_j or momenta $p_j = \sqrt{E_j^2 - m_{\nu_j^2}}$ to arbitrary precision; rather, $\Delta E_j \Delta t_j \gtrsim \hbar/2$ and $\Delta(p_x)_j \Delta x_j \gtrsim \hbar/2$. Correspondingly, the neutrinos actually propagate as wavepackets. As noted above, if the mass differences $|m_{\nu_j} - m_{\nu_k}|$ are sufficiently small relative to the energies $E_j \simeq E_k$, then the resultant velocities are sufficiently close that these wavepackets will continue to maintain a high degree of overlap during the relevant time that they propagate in the experiment, and hence the individual mass eigenstates will remain effectively coherent in the quantum mechanical sense. Their propagation may then be characterized by a single momentum p. Now assume that, after having propagated for a time t (and hence, for $1-v/c \ll 1$, a distance $L = ct$), the neutrino(s) scatter via a charged current weak interaction. This again projects out the weak interaction eigenstates. In particular, because of the different phases which the mass eigenstates pick up during the propagation, a neutrino which is emitted as ν_{ℓ_a} has a nonzero probability to produce a ℓ_b:

$$P = |\langle \nu_{\ell_b}(t) | \nu_{\ell_a}(0) \rangle|^2 = \sin^2(2\theta) \sin^2\Big(\frac{\Delta m^2 L}{4E}\Big) \tag{9}$$

where

$$\Delta m^2 = m_{\nu_i}^2 - m_{\nu_j}^2 \tag{10}$$

Numerically, $\Delta m^2 L/(4E) = (1.266932...)\Delta m^2 L/E$, where Δm^2 is measured in eV^2, L in m, and E in MeV (or L in km, E in GeV). Thus, neutrino oscillation experiments cannot measure individual neutrino masses, but only differences of masses squared, and indeed these are generally weighted in a more complicated way by lepton mixing matrix coefficients for the general case where there is mixing among more than just two species. Experimental results are presented as allowed regions on a plot, the axes of which are $|\Delta m^2|$ and $\sin^2 2\theta$. These are often summarized in terms of the upper limit on Δm^2 (the absolute value is usually suppressed in the notation) for maximal mixing, $\sin^2 2\theta = 1$, and the upper limit on $\sin^2 2\theta$ for "large" Δm^2, *i.e.*, sufficiently large $|\Delta m^2|$ that the detector averages over many cycles of oscillation (or there ceases to be any coherence). A more complete discussion is given in the "Note on Neutrino Oscillation Experiments" just before the tables reporting such results.

An important type of experiment is the search for neutrinoless double-β decay, which tests for total lepton number violation such as would result for Majorana-neutrino masses. This process takes place when a nucleus with Z protons and $A = ZN$ nucleons decays according to $(Z, A) \to (Z+2, A)e^-e^-$, violating total lepton number by two units. In the case of neutrinos with masses which are sufficiently light, an upper limit on neutrinoless double-β decay yields a correlated upper limit on the quantity

$$\overline{m} = \left| \sum_j U_{1j}^2 m_{\nu_j} \right| \ . \tag{11}$$

Cancellations may occur in the sum, since U_{1j} is, in general, complex. The situation is explored further in the minireview by Petr Vogel which prefaces the double-β decay sections. See Ref. 19 for some recent reviews of searches for neutrinoless double-β decay.

A brief summary of the current experimental situation follows (see previous editions for discussions of various positive claims for neutrino masses and mixing, and their refutations).

1. There is no evidence at present from direct searches for nonzero neutrino masses. These include the endpoint of the Kurie plot in nuclear beta decay for $m(\nu_e)$, $\pi^+ \to \mu^+\nu_\mu$ for $m(\nu_\mu)$, and certain τ decays for $m(\nu_\tau)$ (where, as discussed above, the limits actually apply to the respective mass eigenstates ν_1, ν_2, and ν_3 in these three weak eigenstates).

2. There are no indications of any positive neutrino masses from any of the peak search experiments in π or K decay, or from any experiments on neutrino decays.

3. There are no indications of any positive neutrino masses from nuclear beta decay spectra. The 7-year controversy over the claim by Simpson, Hime, and others of a 17 keV neutrino is finally over, with retractions by these authors of their original claims after very strong refutations by a number of high-sensitivity experiments.

4. A number of positive claims for neutrino oscillations in reactor and accelerator neutrino experiments have either been refuted or retracted, or both (see previous editions for details). However, in one analysis, the LSND (Liquid Scintillator Neutrino Detector) group at Los Alamos has reported evidence for the neutrino oscillation $\overline{\nu}_\mu \to \overline{\nu}_e$ [20]. A dissenting analysis of the data by a member of the collaboration reports no evidence for neutrino oscillations [21]. More recently, the LSND group has increased its data sample and strengthened its evidence for neutrino osciallations [22].

5. There is no indication of Majorana neutrino masses from searches for neutrinoless double-β decay.

6. Several experiments have reported evidence for atmospheric neutrino oscillations. Other experiments report results consistent with no such oscillations. This situation is unsettled at present.

7. It is generally acknowledged that the strongest indirect evidence for neutrino masses and mixing is the observed deficit in the solar neutrino flux. The current situation is reviewed by K. Nakamura as a preface to the Solar Neutrino Listings.

For some recent reviews on neutrino physics and further references to the original literature, see Refs. 23–28.

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