NEUTRINOS

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

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 neutrino oscillations and lepton mixing are given in Refs. 1 and 2). Indeed, even in the literature of the 1970's, one will often find statements asserting that in the standard SU(2)× U(1) electroweak theory (without electroweak-singlet neutrinos) the known (electroweak-doublet) neutrinos are massless.

In contrast, in the modern theoretical view based on the standard electroweak theory (and its supersymmetric extensions which stabilize the hierarchy), small but nonzero neutrino masses are expected on general grounds. The reason for this is as follows. Given only the known left-handed neutrino fields and the usual Higgs field(s) in the Standard Model (and supersymmetric extensions thereof which stabilize the hierarchy), nonzero neutrino masses result generically from higher-dimension operators which are strongly expected to occur at a scale near to that of quantum gravity. This expectation is based on the general consensus that pointlike theories of quantum gravity are nonrenormalizable, and is borne out by explicit calculations of the low-energy, field-theory limit of string theories. Such higher-dimension operators would be suppressed by associated inverse powers of the (reduced) Planck mass, $\kappa^{-1} \equiv \overline{M}_P \equiv \sqrt{\hbar c/(8\pi G_N)} = 2.4 \times 10^{18}$ GeV. For example, there would be a gauge-invariant dimension-5 operator consisting of the weak $I = 1$, $Y = -2$ Majorana bilinear ${\nu_L}^T C \nu_L$ (where generation indices are suppressed) contracted with a symmetric, $I = 1$ quadratic product of the usual $I = 1/2$, $Y = 1$ Higgs, with a coefficient of the form $a/\overline{M}_P$, where a is a dimensionless constant. From the vacuum expectation values of the Higgs, this would give rise to (Majorana) neutrino masses of order $m_\nu \sim av^2/\overline{M}_P$, where $v = 250$ GeV is the scale of electroweak-symmetry breaking. (Here $\overline{M}_P$ is an approximate upper bound on the mass which suppresses such operators; it is possible that new physics occurs at some intermediate mass scale $v \ll M_I < \overline{M}_P$ in such a way that dimension-5 operators of this type would give rise to neutrino masses of order av^2/M_I.) This not only leads one to expect nonzero neutrino masses but explains why they are so small.

In retrospect, one sees that this change in theoretical perspective is associated with the change in viewpoint concerning quantum field theory. For decades after the success of quantum electrodynamics in the late 1940's, renormalizability was taken as a necessary property for an acceptable fundamental quantum field theory. It was this implicit theoretical assumption which led to the oft-repeated statements in the 1970's that in the standard electroweak theory (without electroweak-singlet neutrinos) the known neutrinos are massless. More recently, 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 is much smaller than $\overline{M}_P$ (or possibly even some lower scale, M_I), *i.e.*, a consequence of the hierarchy. Once one includes such higher-dimension operators in one's considerations, the realization that neutrino masses are generic follows immediately. A summary of this modern view is given, *e.g.*, in Ref. 3.

In contrast to this mechanism for neutrino mass, which only relies upon the known neutrinos, together with the Higgs field(s) of the Standard Model (or its supersymmetric extensions), there is another more speculative mechanism which we mention here for completeness. It is not known whether there exist any electroweak-singlet neutrino fields. If they do exist, then they could lead, via renormalizable, dimension-4 operators, to neutrino masses $m_\nu \sim v^2/M_R$, where the scale M_R of the electroweak-singlet neutrino mass is naturally $\gg v$, again yielding, albeit for a different reason, very small m_ν [4].

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 [5]. 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 latter are often called "right-handed neutrino singlets," although, since they are singlets, it is a convention whether one chooses to write them as $(\chi_j)_R$ or $(\chi'_j)_L = (\chi^c_j)_L$. 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

$$(\nu_{\ell_a})_L = \sum_j U_{aj}(\nu_j)_L \tag{1}$$

where the $\{\nu_j\}$ denote these mass eigenstates and consist of n members together with possible additional SU(2) $\times$ U(1) singlet neutral leptons, often 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. 6, decays such as ^{3}H $\to$ ^{3}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}H $\to$ ^{3}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 j'th mode in ^{3}H β decay and the k'th 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}H$\to$ ^{3}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 [6,7]. 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 $m_L\,\overline{\nu}_L\,\nu_R^c + h.c. = m_L^*\,\nu_L^{\,T}\,C\nu_L + h.c.$ and right-handed Majorana masses of the form $m_R\,\overline{\chi}_L^c\,\chi_R + h.c. = m_R\,\chi_R^{\,T}C\,\chi_R + h.c.$, where C is the Dirac charge conjugation matrix. 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 coupling $\epsilon_{ij}\widehat{L}^i\widehat{\chi}^c\widehat{H}_u^j$ (all chiral superfields are taken as left-handed), where $\widehat{H}_u$ 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. In supersymmetric extensions of the Standard Model, the χ_R arise as (conjugates of) the spin 1/2 component fields in (left-handed) gauge-singlet chiral superfields. The general neutrino mass term in the Lagrangian is then given by

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

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 [8]. 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 similarity transformation which diagonalizes the neutrino mass matrix, together with the similarity 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 [5,9], as

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

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. 10). Again, however, we note that Dirac neutrinos should be regarded as a special case; the generic case is Majorana neutrinos. The operator products which define the magnetic and electric-dipole moments, *viz.*, $\overline{\nu}\,\sigma_{\alpha\beta}\,\nu\,F^{\alpha\beta}$ and $\overline{\nu}\,\sigma_{\alpha\beta}\,\gamma_5\,\nu\,F^{\alpha\beta}$, respectively (where $F^{\alpha\beta}$ is the electromagnetic field strength tensor) vanish identically if ν is a Majorana neutrino because of the Majorana property that $\nu^c = \pm\nu$. Thus, a Majorana neutrino has identically zero magnetic and electric dipole moments.

Only the diagonal magnetic- and electric-dipole moments are static properties of a given neutrino-mass eigenstate. Transition magnetic and electric dipole moments exist in general for both Dirac and Majorana neutrinos but are not static properties and hence are not considered here. 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. 5 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. 6 was applied to existing data in that paper and a subsequent one [7]; 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

$$\begin{pmatrix} \nu_{\ell_a} \\ \nu_{\ell_b} \end{pmatrix} = \begin{pmatrix} \cos\theta & \sin\theta \\ -\sin\theta & \cos\theta \end{pmatrix} \begin{pmatrix} \nu_i \\ \nu_j \end{pmatrix} \tag{4}$$

where ν_{ℓ_a} are the weak-neutrino eigenstates, with $\nu_{\ell_1} = \nu_e$, *etc.*, and ν_i are neutrino mass eigenstates. Let the distance between the source of the neutrinos and their point of interaction be denoted L, and their energy as E. Assume furthermore that the m_{ν_i} are such that the coherence assumption is valid. Then the probability of an initial ν_{ℓ_a}, having propagated for a distance $t = L$ (with $1 - v/c \ll 1$) being equal to ν_{ℓ_b} is given by

$$P \equiv |\langle \nu_{\ell_b} | \nu_{\ell_a}(t) \rangle|^2 = \sin^2 2\theta \sin^2\Big(\frac{\Delta m^2 L}{4E}\Big) \tag{5}$$

where

$$\Delta m^2 = m^2_{\nu_i} - m^2_{\nu_j} \ . \tag{6}$$

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

$$\overline{m} = \left| \sum_j U_{1j}^2 m_{\nu_j} \right| \ . \tag{7}$$

Cancellations may occur in the sum, since U_{1j} is, in general, complex. See Ref. 11 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-β decay for m_{ν_e}, $\pi^+ \to \mu^+ \nu_\mu$ for m_{ν_μ}, and certain τ decays for m_{ν_τ} (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). There has been some concern over the fact that the quantity which is actually measured, *i.e.*, the square of the neutrino mass, is negative by several standard deviations for ν_e. This may indicate systematic errors in the measurements which are not understood. A similar situation was noticed recently for the muon neutrino mass squared. Although there is still no resolution of this in the refereed journal literature, it has been argued [12,13] that the reason for this is that the charged-pion mass has been determined incorrectly from pionic x-rays, that this determination is actually ambiguous(!), and that when the analysis is interpreted differently, the charged pion mass is increased so as to yield a measurement of $m_{\nu_\mu}^2$ ($m_{\nu_2}^2$) which is no longer tachyonic but is consistent with zero. See the recent preprint [13] and the more extensive discussion before the m_{ν_μ} and $m_{\pi^\pm}$ Listings.

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. These set good upper limits on respective lepton mixing matrix coefficients.

3. There are no indications of any positive neutrino masses from nuclear-β 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. Certain positive claims for neutrino oscillations in reactor and accelerator neutrino experiments have either been refuted or retracted, or both (see previous editions for details).

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

6. The situation concerning possible atmospheric neutrino oscillations remains unsettled.

It is generally acknowledged that the strongest indirect evidence for neutrino masses and mixing is the observed deficit in the solar neutrino flux. The pioneering ^{37}Cl radiochemical solar neutrino experiment of R. Davis and his group started in 1967 and, with only one 18-month gap, has been running ever since. Over the years, a deficit in the measure flux of high-energy neutrinos (mainly from the ^{8}B reaction) has been reported [14,15]. The measured flux is about three times less than the flux predicted by theoretical calculations (Refs. 16–18, and other reviews quoted therein; see also Ref. 10). More recently, Kamiokande, a large water Čerenkov detector, has produced independent high-quality measurements of the solar neutrino flux in roughly the same energy range [19,20,21]. Unlike a radiochemical experiment, a Čerenkov detector preserves directional information, and the Kamiokande collaboration has been able to demonstrate that the neutrinos come from the direction of the sun. They have confirmed the deficiency reported by Davis *et al.* The ratio of observed to expected fluxes seen by the Kamiokande group is somewhat higher than that observed by Davis, but is still significantly below unity.

Two experiments are now operating which are sensitive to the lower-energy neutrinos from the main p-p burning chain. These are SAGE, the Soviet-American Gallium Experiment [22], and GALLEX, a gallium experiment in the Gran Sasso Underground Laboratory [23,24]. Both of these also report a deficiency in the measured solar neutrino flux. These results are especially important since the flux of the *pp* neutrinos constitutes the dominant part of the solar neutrino flux and is widely considered to be reliably calculable. (This is not to imply that calculations of the high-energy neutrino flux are unreliable, only that they are harder to check. The high-energy part does not arise from the reactions producing most of the sun's luminosity, but from minor side-chains which contribute only a very small fraction of the entire solar neutrino flux.)

One explanation of the results of these four experiments is that neutrino oscillations do take place during the transit from the production point in the sun to the interaction point in the earth. These oscillations may involve mixing either with ν_μ and ν_τ or "sterile" neutrino components. An appealing scenario is that of resonant neutrino oscillations, the Mikheyev-Smirnov-Wolfenstein (MSW) effect [25]. Details about this can be found in the reviews cited at the end of this note, and in the references in the original papers.

For some recent reviews on neutrino physics and further references to the original literature, see Refs. [15,26–29].

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Some other sources of information on neutrinos include topical neutrino conferences, and international high-energy physics and lepton-photon conferences.