

MODELS OF NEUTRINO MASSES AND MIXINGS: IMPLICATIONS OF NEUTRINOLESS DOUBLE BETA DECAY

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One of the most important features of neutrinos is that they could be Majorana particles, which violates lepton number. This can be verified in neutrinoless double beta ($0\nu\beta\beta$) decay experiments. An evidence for neutrinoless double beta decay has recently been announced, which has several very strong implications, some of which will be reviewed briefly in this article. If the neutrinos are Majorana particles violating lepton number, then the Majorana masses could explain the baryon asymmetry of the universe. The solar, atmospheric and other laboratory experiments give us the mass squared difference between different neutrinos; only the $0\nu\beta\beta$ decay can give us the absolute masses. The latest result from $0\nu\beta\beta$ decay restricts the patterns of neutrino masses to very few choices, only the degenerate and the inverted hierarchical mass matrices are allowed for three generations of neutrinos. The allowed textures of the neutrino masses are also constrained.

Key Words: Majorana Neutrinos; Lepton Number Violation; Baryon Asymmetry of the Universe; Absolute Mass of Neutrino; Restrictions Imposed by Double Beta Decay

1 Introduction

During the past few years neutrino physics is becoming highly interesting. A small neutrino mass has been evidenced in the atmospheric neutrinos^{1,2} and supported by solar neutrinos^{3,4}. Laboratory experiments are also confirming these results and providing further inputs⁵⁻⁸. Recent result from Kamland⁵ allows only one of the several solutions to the solar neutrino anomaly. Sterile neutrinos are highly disfavoured by all results combined together. With all these advancement of neutrino physics, we can even start thinking of some interesting geophysical and defence applications⁹.

All these experiments provide us with the information about the neutrino oscillations and hence the mass squared differences and the mixing angles between the three neutrinos. However, the absolute values of the masses of the neutrinos are not determined by these experiments. Recently neutrinoless double beta decay has been observed¹⁰, which gives the overall mass scale for the neutrinos. This evidence also tells us that the neutrinos are Majorana particles and hence there is lepton number violation in nature. This

is the first evidence for any lepton number violation in nature. On the other hand, we already know that there are interesting consequences if lepton number is violated. For example, the baryon asymmetry of the universe may be explained naturally if lepton number is broken. The amount of lepton number violation can also constrain some parameters in several models. Since the neutrinoless double beta decay determines the overall mass of the neutrinos, it constrains the possible neutrino mass matrices when combined with the neutrino oscillation results. In this article we shall briefly describe these consequences of the neutrinoless double beta decay.

2 Dirac and Majorana particles

All the quarks and the charged leptons are Dirac particles. Any particles carrying any definite charge are necessarily Dirac particles. But the neutral particles could be either Dirac or Majorana. The main difference between a Dirac and a Majorana particle is that for Majorana particles conservation of all charges is violated (see *e.g.* refs.[10-12]). Since the neutrinos carry only lepton number and no other charges, if they

are Majorana particles then lepton number will be violated by their mass terms.

A Dirac particle satisfies the Dirac equation,

$$i\gamma_\mu \partial_\mu \psi_D - m_D \psi_D = 0, \quad \dots(1)$$

which follows from the Lagrangian, $\mathcal{L} = \mathcal{L}_{KE} + \mathcal{L}_D$ with

$$\begin{aligned} \mathcal{L}_{KE} &= i\bar{\psi}_D \gamma_\mu \partial_\mu \psi_D \\ &= i\bar{\psi}_L \gamma_\mu \partial_\mu \psi_L + i\bar{\psi}_R \gamma_\mu \partial_\mu \psi_R, \quad \dots(2) \end{aligned}$$

$$\begin{aligned} \text{and } \mathcal{L}_D &= -m_D \bar{\psi}_D \psi_D = -m_D \bar{\psi}_R \psi_L + h.c. \\ &= m_D \psi_L^T C^{-1} \psi^c_L + h.c. \quad \dots(3) \end{aligned}$$

where $h.c.$ stands for Hermitian conjugate term. ψ_L and ψ_R are the projections of the Dirac spinors ψ_D into its left-handed and right-handed Weyl components, so that $\psi_D = \psi_L + \psi_R$, with

$$\psi_L = \frac{1}{2}(1 - \gamma_5)\psi_D \quad \text{and} \quad \psi_R = \frac{1}{2}(1 + \gamma_5)\psi_D, \quad \dots(4)$$

which satisfies, $\gamma_5 \psi_L = -\psi_L$ and $\gamma_5 \psi_R = \psi_R$. The kinetic energy term thus treats the left-handed and the right-handed fields independently, while the mass term connects the left-handed fields to the right-handed fields.

The Dirac mass term can also be written in a more general way as

$$\begin{aligned} \mathcal{L}_D &= \frac{1}{2}(\psi_L \quad \psi^c_L)^T C^{-1} \begin{pmatrix} 0 & m_D \\ m_D & 0 \end{pmatrix} \begin{pmatrix} \psi_L \\ \psi^c_L \end{pmatrix} \\ &= \frac{1}{2}\Psi_L^T C^{-1} M_D \Psi_L, \quad \dots(5) \end{aligned}$$

where M_D is the 2×2 Dirac mass matrix and $\Psi_L^T = (\psi \quad \psi^c)_L^T$.

A Majorana particle is its own anti-particle

$$\psi_M^c = \eta_C \psi_M \quad \dots(6)$$

where η_C is a new phase, called a creation phase or a Majorana phase. This Majorana condition relates the creation and annihilation operators for the particles and antiparticles, and hence there is no vector current for any Majorana particle. The number operator also vanishes, $\int d^3x \bar{\psi}_M \gamma_0 \psi_M = 0$ for the Majorana particles, so there is no conservation of number of particles. If any Majorana particle carries some

U(1) charge (say, the lepton number), this charge (the lepton number) is violated since the particles and the antiparticles carry opposite charges. In this case one can write down a mass term which violates the U(1) charge by two units

$$\mathcal{L}_M = -\frac{1}{2}m_M \bar{\psi}_M^c \psi_M = \frac{1}{2}m_M \psi_M^T C^{-1} \psi_M. \quad \dots(7)$$

Unlike the Dirac mass, the Majorana mass term does not mix the left-handed particles with the right-handed particles. The Majorana mass term takes a left-handed particle to its $\mathcal{C}\mathcal{P}$ conjugate, the right-handed anti-particle. Thus a Majorana particle can be written in the Weyl basis as $\psi_M = \psi_L + \eta_C^* \psi_R^c$.

Since the Majorana mass connects a left-handed particle with its $\mathcal{C}\mathcal{P}$ conjugate, and similarly for a right-handed particle, it is possible to write down the Majorana mass terms for the left-handed and right-handed particles independently,

$$\begin{aligned} \mathcal{L}_L &= -\frac{1}{2}m_L \bar{\psi}_L^c \psi_L = \frac{1}{2}m_L \psi_L^T C^{-1} \psi_L \\ &= -\frac{1}{2}m_L \bar{\psi}_R \psi_R^c \\ \mathcal{L}_R &= -\frac{1}{2}m_R \bar{\psi}_R^c \psi_R = \frac{1}{2}m_R \psi_R^T C^{-1} \psi_R \\ &= -\frac{1}{2}m_R \bar{\psi}_L \psi_L^c. \quad \dots(8) \end{aligned}$$

Thus the left-handed particles and the right-handed particles can be treated as two independent Majorana particles,

$$\psi_{M1} = \psi_L + \eta_C^* \psi_R^c \quad \text{and} \quad \psi_{M2} = \psi_R + \eta_C^* \psi_L^c, \quad \dots(9)$$

with Majorana masses $M_1 = m_L$ and $M_2 = m_R$.

When there are both left-handed and right-handed neutrinos in any theory, it is convenient to write down the Majorana masses in the basis $\Psi_L^T = (\psi \quad \psi^c)_L^T$,

$$\begin{aligned} \mathcal{L}_M &= \frac{1}{2}(\psi \quad \psi^c)_L^T C^{-1} \begin{pmatrix} m_L & 0 \\ 0 & m_R \end{pmatrix} \begin{pmatrix} \psi \\ \psi^c \end{pmatrix}_L \\ &= \frac{1}{2}\Psi_L^T C^{-1} M_M \Psi_L, \quad \dots(10) \end{aligned}$$

where M_M is the 2×2 Majorana mass matrix. Thus in this basis the only difference between the Dirac and the Majorana mass term is that the Majorana mass matrix is diagonal, while the Dirac mass matrix contains only the off-diagonal terms.

For neutrinos, when both left-handed and right-handed neutrinos are present, the mass matrix will

have two components $M_V = M_D + M_M$, where the Dirac part contains the off-diagonal terms and the Majorana part contains non-vanishing diagonal elements. From now on we simplify our notation and write the mass matrix as,

$$\mathcal{L}_M = (\psi \quad \psi^c) \begin{pmatrix} m_L & m_D \\ m_D & m_R \end{pmatrix} \begin{pmatrix} \psi \\ \psi^c \end{pmatrix} \dots (11)$$

In this case to find out the nature of the physical particles one has to diagonalize the mass matrix and then identify the physical states. The physical particles could now be¹²

- (i) *Dirac particle*: If $M_M = 0$, the two states ψ_{M1} and ψ_{M2} will combine to form a single Dirac particle. Lepton number will be conserved in this case.
- (ii) *pseudo-Dirac particle*: If the two eigenvalues are the same with opposite sign, but $M_M \neq 0$, the two states will combine to form a Dirac particle at the tree level. But the radiative corrections will split the degeneracy leading to two Majorana particles with almost equal masses. Lepton number will be broken in this case, and the amount of lepton number violation will depend on the tree level masses M_M .
- (iii) *Majorana particles*: If the two eigenvalues are different, this will represent two physical Majorana particles. Lepton number will be violated in this case.

Thus $M_M \neq 0$ is always associated with lepton number violation. The physical particle could then be a pseudo-Dirac or two Majorana particles.

3 Neutrinoless Double Beta Decay

A direct consequence of the lepton number violation is neutrinoless double beta decay (see e.g., ref.[13]). In some even-even nuclei ordinary beta decay is forbidden because of the masses of the nuclei. However, double beta decay could still be allowed, including neutrinoless double beta decay

$$n + n \rightarrow p + p + e^- + e^- + \nu_e + \nu_e \dots (12)$$

$$n + n \rightarrow p + p + e^- + e^- \dots (13)$$

There are about 35 candidates for neutrinoless double beta decay in nature. For a few of them the two-neutrino decay mode has been observed experimentally (see ref.[13]), while the neutrinoless decay mode

has not been observed before the papers in ref.[10]. The neutrinoless double beta decay (eq.13) process would violate lepton number by two units ($\Delta L = 2$), and directly measures the amount of lepton number violation. In the $2\nu\beta\beta$ decay (eq.12) the total kinetic energy of the two electrons has a distribution, since part of the kinetic energy is carried by the neutrinos. But for the $0\nu\beta\beta$ decay the total kinetic energy of the two electrons is constant and is given by the Q value of the decay. Thus if the energy resolution is high enough, it is possible to measure the $0\nu\beta\beta$ decay mode without any background from the $2\nu\beta\beta$ decay.

Let us first make the most natural assumption, namely that the $0\nu\beta\beta$ process is triggered dominantly by the neutrino mass mechanism, *i.e.*, by exchange of a Majorana neutrino (see Fig.1). In the $0\nu\beta\beta$ decay, only the electron neutrino (ν_e , which has charge current interaction with the electrons) enters in the diagram. The physical states (the states with definite mass) are not the ones which appear in the interaction with the weak gauge bosons.

The neutrinos with definite flavor ($\nu_{\alpha L}$, which has charge current interaction with the charged leptons $\alpha = e^-, \mu^-, \tau^-$) are related to the mass eigenstates or the physical states (ν_{iL} , with masses m_i) through the unitary mixing matrix $U_{\alpha i}$ by

$$\nu_{\alpha L} = \sum_{i=1}^3 U_{\alpha i} \nu_{iL} \dots (14)$$

Due to this mixing lepton flavours (L_e, L_μ, L_τ) are not conserved. However, the total lepton number ($L = L_e + L_\mu + L_\tau$) conservation depends on whether the neutrinos are Dirac or Majorana particles. The mass matrix in the flavor basis, *i.e.*, the basis in which the charged lepton mass matrix is diagonal, is then given by,

$$M_{\alpha\beta}^V = U_{\alpha i} M_{ij}^{\text{diag}} U_{\beta j}^T \dots (15)$$

where $M_{ij}^{\text{diag}} = m_i \delta_{ij}$.

When the neutrinos are Majorana particles there will be total lepton number L violation, which will allow neutrinoless double beta decay through the diagram given in Fig.1. Here the neutrinos are virtual particles in the intermediate state, so the neutrino masses and mixing enter into the neutrino propagator. Since the physical electrons are emitted in the process, the virtual neutrinos are electron flavor neutrinos and

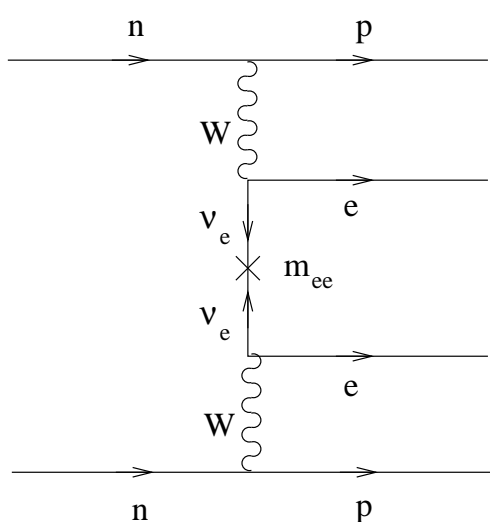


Fig. 1 Diagram for neutrinoless double beta decay

hence the effective mass that enters in the process is¹⁴

$$\langle m \rangle = \sum_{i=1}^3 U_{ei}^2 m_i = m_{ee} = M_{ee}^V \quad \dots (16)$$

where m_{ee} is the (11) element of the neutrino mass matrix when written in the flavor basis. Thus the neutrinoless double beta decay measures the lepton number violation and hence the Majorana mass of the neutrinos directly.

The amplitude for the neutrinoless double beta decay, which is derived from the half life measurement is given by¹⁵

$$A_{(\beta\beta)_{0\nu}} = \frac{G_F^2}{2\pi^3} \sum_{i=1}^3 U_{ei}^2 m_i F(m_i, A) \quad \dots (17)$$

where^{13, 15}

$$F(m_i, A) = \frac{\langle e^{-m_i r}/r \rangle}{\langle 1/r \rangle} \quad \dots (18)$$

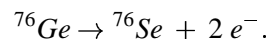
is the nuclear factor, r is the distance between the two neutrons within the nucleus, $R \geq r/2$ is the radius of the initial nucleus and the average is with respect to the two nucleon correlation function appropriate to a nucleus of atomic number A . $F(m_i, A)$ can be approximated by a step function,

$$\begin{aligned} F(m_i, A) &= 1 & \text{for } m_i \ll 10 \text{ MeV} \\ F(m_i, A) &= 0 & \text{for } m_i \gg 10 \text{ MeV.} \end{aligned} \quad \dots (19)$$

Thus here it is assumed that heavier states will not contribute to the neutrinoless double beta decay.

The Heidelberg-Moscow $0\nu\beta\beta$ decay experiment in the Gran Sasso underground laboratory operates

five enriched (to 86%) high purity ^{76}Ge detectors, with a total mass of 11.5 kg, the active mass of 10.96 kg being equivalent to a source strength of 125.5 mol ^{76}Ge nuclei. They look for the decay mode



The high resolution of the Ge detectors eliminates the $2\nu\beta\beta$ decay background. The Q value for this decay is known with high precision¹⁶ to be $Q_{\beta\beta} = 2039.006(50)$ keV, which is very useful in identifying the experimental signature of the $0\nu\beta\beta$ decay mode, which is a peak at the Q -value of the decay.

The data taken during the period August 1990 – May 2000 has been analyzed, also using refined digital pulse shape analysis technique in the energy interval 2000 - 2080 keV around the Q value. Evidence for the $0\nu\beta\beta$ decay mode has recently been announced¹⁰. With the Bayesian method the evidence for this decay mode is at the 97.4 % c.l. (2.2σ) and using the Feldman-Cousins method, also recommended by the Particle Data Group it is more than 99.8 % c.l. (3.1σ). The half-life for the $0\nu\beta\beta$ decay is found to be

$$T_{1/2}^{0\nu} = (0.7 - 18.3) \times 10^{25} \text{ y} \quad (95 \% \text{ c.l.})$$

with a best value of 1.5×10^{25} y. The present signal for the neutrinoless double beta decay amounts to an effective Majorana neutrino mass of the electron neutrino in the range of¹⁰

$$\langle m \rangle = (0.05 - 0.84) \text{ eV} \quad (95\% \text{ c.l.}) \quad \dots (20)$$

with a best value of 0.39 eV. This value allows for a $\pm 50\%$ uncertainty in the nuclear matrix element. This number is in agreement with the limits obtained by all other double beta decay experiments¹⁷ which reach continuously less sensitively (for details see ref.[10]).

4 Constraints on Lepton Number Violation

The claimed evidence and the upper limit also set strong constraints on any lepton number violating processes that may contribute to the $0\nu\beta\beta$ decay diagram, the most conservative ones arising when assuming that the neutrino mass mechanism gives only a minor contribution to the amplitude of the $0\nu\beta\beta$ decay.

The heavy right-handed Majorana neutrinos have small mixing with the left-handed neutrinos, which

allow these heavy neutrinos to enter in the $0\nu\beta\beta$ decay. The upper limit on the lifetime of the $0\nu\beta\beta$ decay then gives a lower limit on the mass of the heavy neutrinos¹⁸

$$M_N > 9 \times 10^7 \text{ GeV}. \quad \dots(21)$$

This implies a lower bound on the mass eigenstates of the heavy neutrinos to be $M_{N_i} > 5 \times 10^5 \text{ GeV}$, with an assumption on the mixing matrix.

Using the lifetime of $0\nu\beta\beta$ decay, the probability for the discovery of the inverse beta decay process $e^-e^- \rightarrow W^-W^-$ at NLC could be constrained. The present value¹⁰ can be achieved at a future linear collider NLC when it reaches a center of mass energy¹⁹ of 2 TeV. The discovery limit for $e^-e^- \rightarrow W^-W^-$ at the NLC for several center-of-mass energies is presented in Fig.2 as a function of the heavy neutrino mass eigenstate M_i and their mixing with lighter eigenstates with masses m_α ($\alpha = e, \mu, \tau$), U_{ei}^2 . This figure considers 10 events as discovery limit and assumes an unpolarized e^- beam with a luminosity of $80 [\sqrt{s}(1 \text{ TeV})]^2 \text{ fb}^2$. From the figure it is clear the for $\sqrt{s} = 500 \text{ GeV}$ and 1 TeV, the parameter space which produces observable $e^-e^- \rightarrow W^-W^-$ is already ruled out by the present result on the $0\nu\beta\beta$ decay. Other lepton number violating processes like $e^- \gamma \rightarrow e^+W^-W^-$ are even less probable and may be observed at NLC only above $\sqrt{s} > 4 \text{ TeV}$.

Some other extensions of the standard model are also constrained by the present limit from $0\nu\beta\beta$ decay. The bound on the the right-handed W boson is¹⁸

$$m_{W_R} \geq 1.2 \left(\frac{M_N}{1 \text{ TeV}} \right)^{-1/4} \text{ TeV}. \quad \dots(22)$$

With some theoretical input, this may be translated to an absolute lower bound of $m_{W_R} > 1.2 \text{ TeV}$ ¹⁸.

If there are Higgs scalar bilinears, which couples to the usual quarks and leptons, they can also allow for $0\nu\beta\beta$ decay. The dileptons were first considered²⁰ (see Fig.3) in connection with the left-right symmetric model, but the bounds are very weak in this class of diagrams.

For the leptoquarks, the bound is even stronger²¹. If the X -type leptoquarks ($SU(2)_L$ singlets) mix with the Y -type leptoquarks ($SU(2)_L$ doublets), then they can give an effective operator $u\bar{\nu}d\bar{l}$. This can then generate a diagram contributing to the neutrinoless double beta decay involving the leptoquarks (Fig.4). A mixing between these two leptoquarks could take

place only after the electroweak symmetry breaking, if both these leptoquarks couple to the usual standard model Higgs doublet ϕ . In that case a coupling ϕXY will induce a mixing of X with Y when ϕ acquires a vev .

The couplings of ϕ with the leptoquarks have been studied²¹ and the present lifetime for the neutrinoless double beta decay is used to constrain these couplings and the leptoquark masses. It was noticed²¹ that in the leptoquark mediated case, there is a huge enhancement factor of $\langle q \rangle / m_\nu \sim 10^8$ ($1\text{eV}/m_\nu$), where m_ν is the effective neutrino mass entering the neutrinoless double beta decay contribution, and $\langle q \rangle$ is the Fermi momentum of a nucleon inside a nucleus, which is about 200 – 300 MeV. For a leptoquark with mass of the order of 100 GeV, the effective coupling constant (including the mixing contribution) comes out to be about 10^{-9} . Thus even if there are any light leptoquarks, they should not be observed in the near future.

There are other types of exotic scalar bilinear mediated diagrams contributing to the neutrinoless double beta decay (Fig.5), which can also constrain the masses and couplings of these exotics²². Assuming, for simplicity, a common mass for all these scalars of the order of about 100 GeV it is possible to obtain a strong constraint on the couplings of all the scalars with the usual fermions. Assuming further the self interactions of the Higgs to be of the order of 1, the bound on the effective coupling of the scalars becomes $f < 10^{-7}$. This can be translated to a bound on the scalar mass assuming all coupling constants to be of the order of 1. A generic bound on the masses of the scalars is then given by the neutrinoless double beta decay, which is of the order of 10^{-4} GeV . Even though this bound is not too strong, this rules out the possibility of detecting several of the scalars in the next generation accelerators, when combined with other constraints.

Considering all possible diagrams discussed above, it is possible to constrain the ratio of the masses to their coupling to first generation fermions for all the exotic scalar bilinears from the neutrinoless double beta decay lifetime. Some of these bounds are comparable to the bounds from other processes like $K^\circ - \bar{K}^\circ$ oscillations, $B^\circ - \bar{B}^\circ$ mixing, $D^\circ - \bar{D}^\circ$ mixing, proton decay or $n - \bar{n}$ oscillations, all of these bounds are much stronger than the direct bound from the accelerators.

The indirect bounds discussed also constrain some

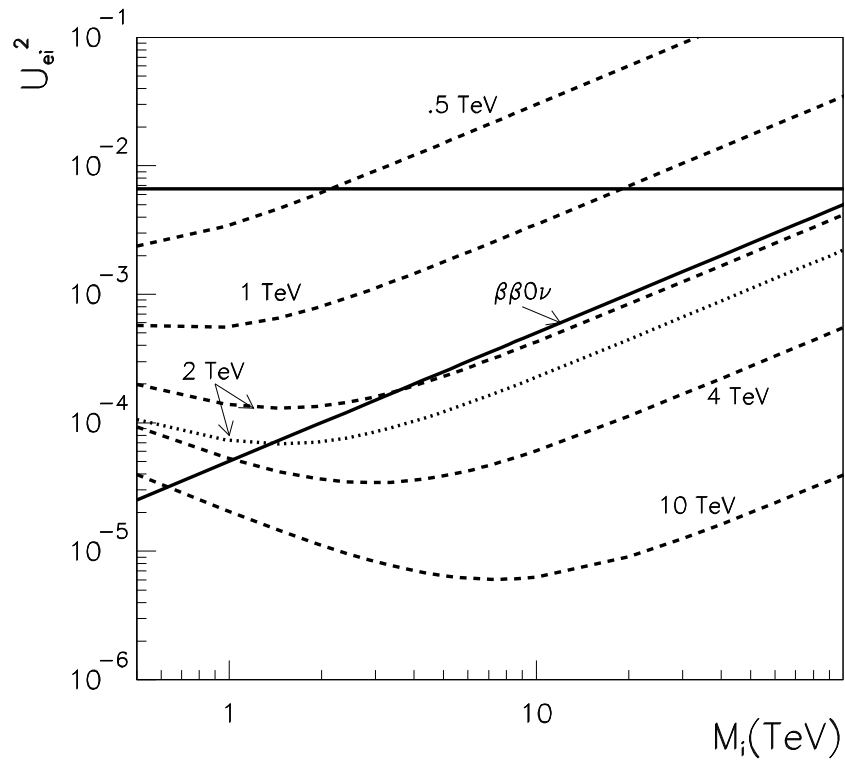


Fig. 2 Discovery limit for $e^-e^- \rightarrow W^-W^-$ at the NLC is compared with the $0\nu\beta\beta$ decay result^{10,17} (from ref.[19])

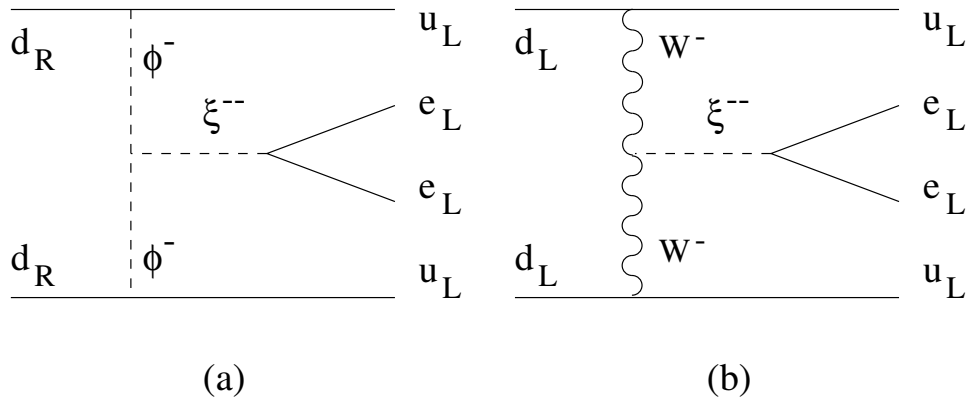


Fig. 3 Diagram for neutrinoless double beta decay with dileptons.

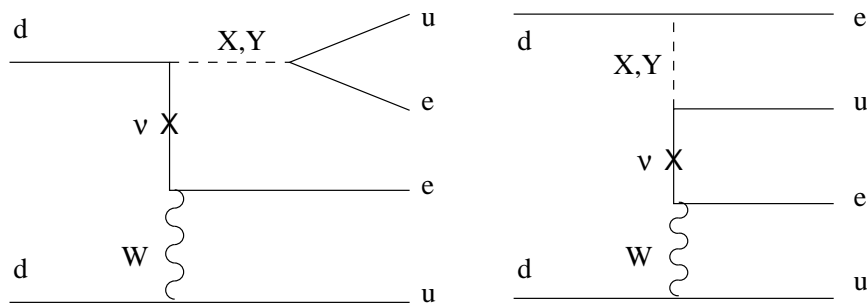


Fig. 4 Diagram for neutrinoless double beta decay with leptoquarks (from ref.[21])

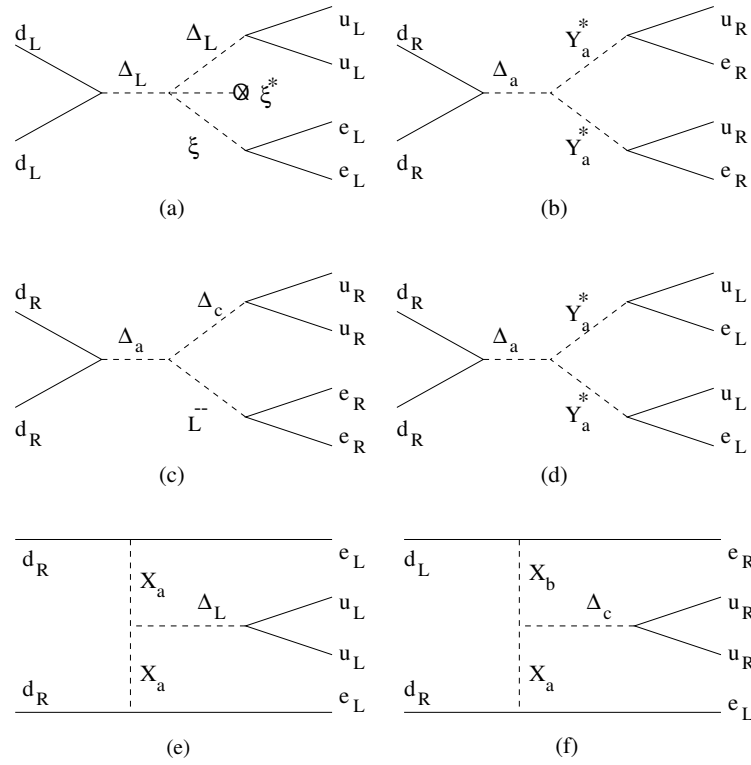


Fig. 5 Diagram for neutrinoless double beta decay with scalar bilinears

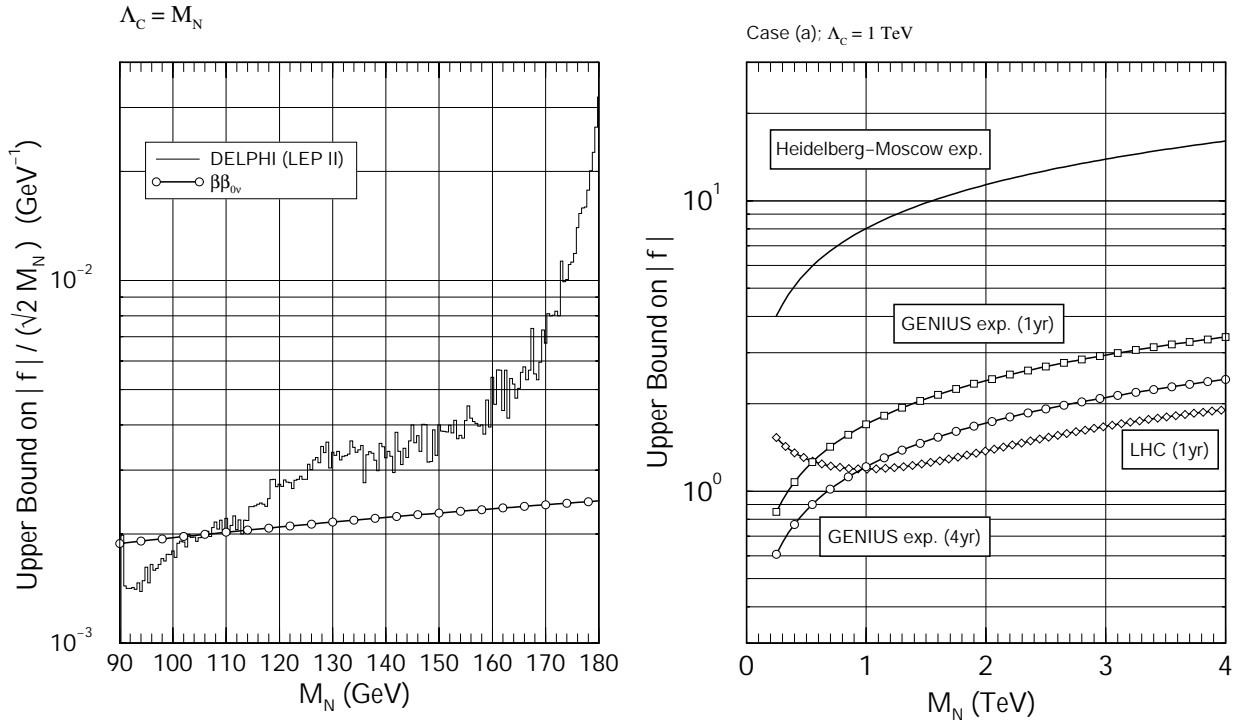


Fig. 6 (a) Comparison between the $\beta\beta_{0\nu}$ and the LEP II upper bound on $|f|/(\sqrt{2}M_N)$ as a function of M_N for $\Lambda_c = M_N$. Regions above the curves are excluded. (b) Comparison between the $\beta\beta_{0\nu}$ result from GENIUS and the LHC upper bound on $|f|/(\sqrt{2}M_N)$ as a function of M_N for $\Lambda_c = M_N$. Regions above the curves are excluded (from ref.[23])

of the possibilities of composite particles. For example, if a neutrino is a composite particle, the most severe constraint comes from the neutrinoless double beta decay²³,

$$|f| \leq 3.9 \frac{\Lambda_c}{1 \text{ TeV}} \left(\frac{M_N}{1 \text{ TeV}} \right)^{1/2} \dots (23)$$

where Λ_c is the compositeness scale, M_N is the mass of the heavy excited neutrino and f is the dimensionless coupling constant. A comparison of the bound on the ratio of the mass to coupling, from LEP II and the neutrinoless double beta decay is shown in ref.[23]. It has been shown by ref.[23], that the limit from the Heidelberg-Moscow experiment is already stronger than that given by LEP II (Fig.6a). In the next generation neutrinoless double beta decay experiments, the constraint may be further improved. The proposed sensitivity by the GENIUS may be able to give a bound, which will be comparable to the LHC bound, as shown in Fig.6b.

It is also possible to constrain several parameters of the supersymmetric theories. All sources of lepton number violation could contribute to the $0\nu\beta\beta$ decay and hence they are all constrained. In R-parity conserving theories, any mechanism of Majorana neutrino masses will imply lepton number violation and the amount of constraint will be given by the present value of the neutrino masses. In addition, there is a new source of lepton number violation coming from soft supersymmetry breaking, which gives rise to sneutrino mixing²⁴. The amount of sneutrino-antisneutrino oscillation is also constrained by the present limit on the $0\nu\beta\beta$ decay²⁴.

The mass difference between a sneutrino and an antisneutrino, \tilde{m}_M , comes from lepton number violation. The $0\nu\beta\beta$ decay gives a bound on the effective lepton number violating MSSM parameter $\eta^{susy} < 1 \times 10^{-8} (m_{susy}/100 \text{ GeV})^5$. Assuming all superparticles to have a common mass m_{susy} , this amounts to a bound²⁴

$$\tilde{m}_M < 2 \left(\frac{m_{susy}}{100 \text{ GeV}} \right)^{3/2} \text{ GeV}, \dots (24)$$

when the neutralino is mostly a Bino, and

$$\tilde{m}_M < 11 \left(\frac{m_{susy}}{100 \text{ GeV}} \right)^{7/2} \text{ GeV}, \dots (25)$$

when the neutralino is mostly a Higgsino. There are also bounds on \tilde{m}_M coming directly from the neutrino

masses with some assumptions, which are $\tilde{m}_{M(e)} < 120 \text{ MeV}$, $\tilde{m}_{M(\mu)} < 13 \text{ GeV}$ and $\tilde{m}_{M(\tau)} < 149 \text{ GeV}$. Thus for the second and third generations large mixing is consistent with present limits.

In R-parity violating theories one may assume that the baryon number violating couplings are small, while lepton number violating couplings are unconstrained by theory. In this case, all the lepton number violating couplings can give rise to new contributions to the neutrinoless double beta decay²⁵. The present result on the half-life¹⁰ of $0\nu\beta\beta$ decay then gives a strong bound for the first generation

$$\lambda'_{111} < 3.3 \times 10^{-4} \left(\frac{m_{\tilde{q}}}{100 \text{ GeV}} \right)^2 \left(\frac{m_{\tilde{g}}}{100 \text{ GeV}} \right)^{1/2}, \dots (26)$$

assuming equal mass for the up-squark and the down-squark. An example of the type of diagrams contributing to the $0\nu\beta\beta$ decay, where lepton number violation comes only from R-parity violating couplings, is given in Fig.7a. The above bound comes from such diagrams. There are also diagrams (as shown in Fig.7b), where lepton number violation comes from R-parity violation as well as neutrino mass²⁵. Such diagrams give bounds on products of two couplings,

$$\begin{aligned} \lambda'_{112} \lambda'_{121} &< 1.1 \times 10^{-6} \left(\frac{\Lambda_{susy}}{100 \text{ GeV}} \right)^3 \\ \lambda'_{113} \lambda'_{131} &< 3.8 \times 10^{-8} \left(\frac{\Lambda_{susy}}{100 \text{ GeV}} \right)^3 \dots (27) \end{aligned}$$

for supersymmetric mass parameters of the order of 100 GeV.

5 Constraints on Neutrino Masses

The Heidelberg-Moscow experiment gave¹⁰ the first evidence for lepton number violation in nature. Several implications of this result have been worked out in the literature^{26, 27, 29}. We highlight here a few salient features of this result. The positive indication of neutrinoless double beta decay means that the neutrinos are Majorana particles and hence all the models of Dirac neutrino masses become redundant. Two Majorana neutrinos with opposite $\mathcal{C}\mathcal{P}$ phase and equal mass can form a Dirac particle, canceling their contribution to the neutrinoless double beta decay, that possibility is also ruled out. The smallness of the Majorana mass implies a large lepton number violating

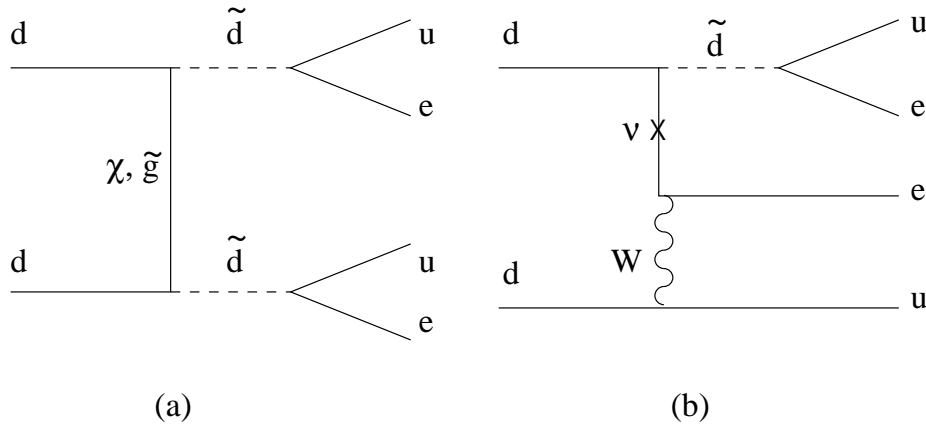


Fig. 7 R-parity violating terms contributing to neutrinoless double beta decay

scale, which can then generate a lepton asymmetry of the universe and explain why there is more matter in the universe compared to anti-matter.

Now the most interesting scenario is the three generation model of Majorana neutrinos. Even in this case it is possible that in some models the contributions of the three neutrinos to the neutrinoless double beta decay cancel each other, which are now eliminated. This forbids some of the structures of the neutrino mass matrix. Several of the the possible solutions of the neutrino masses and mixing, allowed by the solar and atmospheric neutrinos are ruled out, when the neutrinoless double beta decay results are included^{26,27}. With the neutrinoless double beta decay result the electroweak radiative corrections become important in some cases, which then rule out the VAC solutions of the solar neutrinos²⁷. The VAC solution is also ruled out by the recent result from Kamland²⁸.

In the three generation case the contribution to the neutrinoless double beta decay is

$$m_{ee} = \sum_i |U_{ei}|^2 m_i = m_1 c_1^2 c_3^2 + m_2 s_1^2 c_3^2 + m_3 s_3^2. \quad \dots(28)$$

m_i can now be complex and the Majorana phases can contribute to the $0\nu\beta\beta$ decay, except for the overall phase. Considering the limit on s_3 to be given by the CHOOZ result and s_1 from solar neutrino measurements, for each of the possible mass patterns it is possible to examine the neutrinoless double beta decay contributions^{26,27}.

The hierarchical scenario allows a very small m_1 and with all the allowed values of s_1 and s_3 the contribution to m_{ee} comes out to be smaller than the allowed limit for the SMA, LMA and also the VAC solutions of the solar neutrinos. Another possibility with hierar-

chical mass matrix (triple) is to consider all elements of the mixing matrix to be equal, which has a definite prediction for the $0\nu\beta\beta$ decay $m_{ee} \sim 0.02$ and is also ruled out. Hence the neutrinoless double beta decay result rules out the hierarchical patterns of neutrino mass matrices for all the solutions of the solar neutrinos. This is true for the partial degenerate scenarios as well, in which case the constraint on the neutrino mass coming from the atmospheric neutrinos, $0.005 \text{ eV} < m_1 < 0.042 \text{ eV}$, restricts the amount of $0\nu\beta\beta$ decay to be less than the allowed range. Only the degenerate and the inverted hierarchical scenarios are allowed by the present result. This is shown in the summary Fig.8.

In the degenerate case the main contribution to the neutrinoless double beta decay comes from the m_1 , which is greater than the atmospheric neutrino solution^{26,27}, $0.042 \text{ eV} \leq m_1 \leq 1 \text{ eV}$. For both LMA and SMA solutions of the solar neutrinos the contribution could be large and within the range evidenced recently. The LMA and SMA solutions are also allowed for the inverted hierarchical scenario, where the main contribution to the $0\nu\beta\beta$ decay comes from the heavier ν_e state, whose mass is given by the atmospheric neutrinos, so that $m_{ee} < .063$.

In both these cases the VAC solutions to the solar neutrinos are unstable against radiative corrections and are ruled out²⁷. Consider a neutrino mass matrix in the flavor basis. The (11) element of the mass matrix gives the amount of $0\nu\beta\beta$ decay. For the degenerate or the inverted hierarchical scenario either the (22) or the (33) element has to be approximately equal to the (11) element. It is possible to have (23) and (32) elements equal to twice the (11) element, but then the maximal mixing required by the VAC solutions

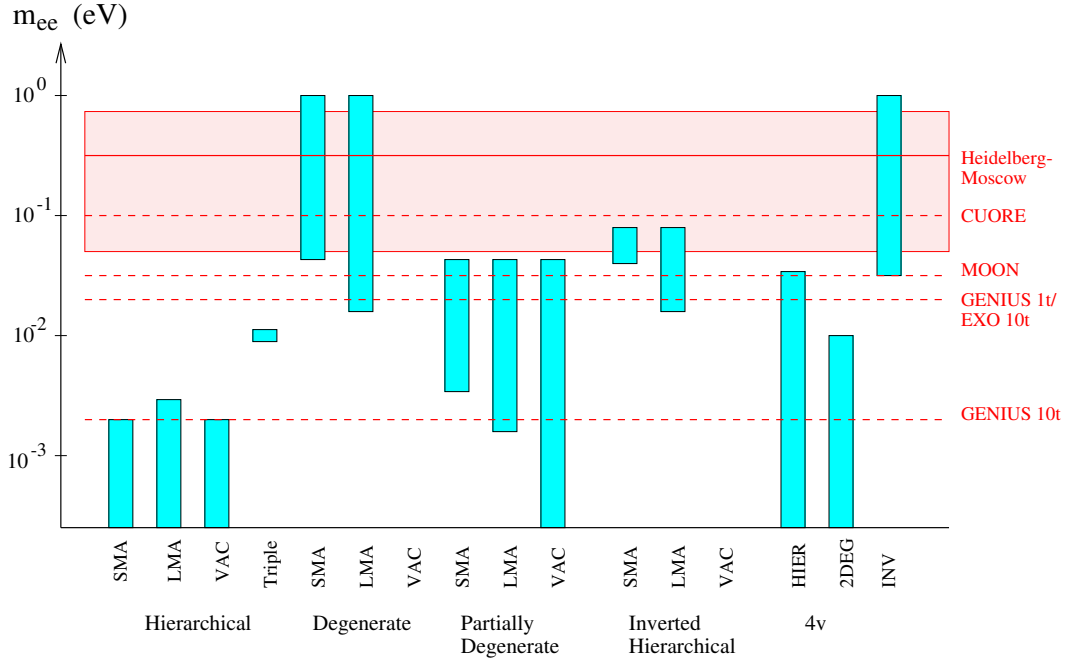


Fig. 8 Contributions in different models to the neutrinoless double beta decay. The present result is given by the shaded region (the solid line denoting the best value), which allows only few of the models. The expected sensitivity for the CUORE³⁰, MOON³¹ and the one ton and ten tons GENIUS³² are given for comparison (from ref.[27])

would give a large mass-squared difference. So, the neutrinoless double beta decay, maximal mixing for the VAC solution and the approximate degeneracy together imply that the (33) element has to be approximately equal to the (11) element, which is greater than the mass required by the atmospheric neutrinos $m_0 > .042$ eV. Then the radiative corrections would introduce a mass-squared difference

$$\begin{aligned} \Delta m^2 &\sim \alpha_w m_0^2 \left(\frac{m_\tau^2}{m_W^2} \right) \sim 1.7 \times 10^{-5} m_0^2 \text{ eV}^2 \\ &> 1.5 \times 10^{-8} \text{ eV}^2 \dots (29) \end{aligned}$$

α_w is the weak fine structure constant. With this correction it will not be possible to maintain the mass squared difference required for the VAC solutions. The VAC solutions are thus ruled out for both the degenerate and inverted hierarchical scenarios and hence for all the three generation scenarios by neutrinoless double beta decay result. The VAC solution is also ruled out by the recent result from KamLAND⁵.

The small mixing angle MSW solution of the solar neutrino problem is also ruled out by the results from SNO and KamLAND. An earlier analysis had pointed out³³ that the requirement of dark matter and SMA MSW solution would not be consistent with Majorana

neutrinos, or the neutrinoless double beta decay is not consistent with SMA MSW solution and dark matter.

There are cosmological restrictions on the number of neutrino flavours. Although three generation scenario is favored, a fourth sterile neutrino is not ruled out. For such four generation scenario (three active and one sterile neutrinos, all of which are Majorana particles), the neutrinoless double beta decay experiment allows only the inverted hierarchy type mass matrix (see Fig.8). At present this solution seems to be highly improbable. When SNO result is combined with other solar neutrino results, an oscillation of ν_e to a sterile neutrino cannot explain the solar neutrino problem. The reactor anti-neutrino disappearance at KamLAND also disfavours sterile neutrinos for the solution of the solar neutrino. The atmospheric neutrino also cannot involve a sterile neutrino. Thus all these results now point towards a restricted form of a three generation neutrino mass matrix.

6 Neutrino Mass Textures

The maximal mixing required for the solutions of the atmospheric neutrinos allows only certain textures of neutrino mass matrices in the three generation scenario^{34,35}. These texture mass matrices are the ones

with few non-vanishing large entries, to which one can add small perturbations to form the complete neutrino mass matrix. There are models in which such allowed texture mass matrices have been incorporated³⁶.

In the flavor basis, in which the charged lepton mass matrix is diagonal, there are only three possible texture mass matrices which can explain the $0\nu\beta\beta$ decay result and are consistent with other experiments³⁷. These are

$$M_V^{A1} = m_0 \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1/2 & 1/2 \\ 0 & 1/2 & 1/2 \end{pmatrix} \quad M_V^{A2} = m_0 \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 1 \end{pmatrix}$$

$$M_V^{A3} = m_0 \begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & 1 & 0 \end{pmatrix}, \quad \dots(30)$$

m_0 is the overall scale in these mass matrices, which is determined by the value of the neutrinoless double beta decay.

There are only a few possible perturbations possible for each of these textures. We present some simple perturbations below (see ref.[37]):

$$m_1^{A1} = m_0 \begin{pmatrix} a & b_1 & b_2 \\ b_1 & 0 & 0 \\ b_2 & 0 & 0 \end{pmatrix}. \quad \dots(31)$$

In this case it is possible to have solutions with $m_0 = 0.05$ eV; $a = 0.003$ and $b_1 = b_2 = 0.001$. It is also possible to have solutions with $b_1 = 0$ and $b_2 = b$ or $b_1 = b$ and $b_2 = 0$, with $b = 0.002$. The last two solutions are equivalent. These are the partially degenerate solutions where two of the masses are degenerate. Then the contribution to the $0\nu\beta\beta$ decay is restricted by the mass required for the atmospheric neutrinos, since $\Delta m_{am}^2 = m_0^2$.

Few possible perturbations with the other texture mass matrices are

$$m_1^{A2} = m_0 \begin{pmatrix} a & b & b \\ b & 0 & b \\ b & b & 0 \end{pmatrix}$$

and $m_1^{A3} = m_0 \begin{pmatrix} a & 0 & a \\ 0 & 0 & 0 \\ a & 0 & -2a \end{pmatrix}.$

$$m_2^{A3} = m_0 \begin{pmatrix} a & a & a \\ a & 0 & 0 \\ a & 0 & -2a \end{pmatrix}$$

and $m_3^{A3} = m_0 \begin{pmatrix} a & 0 & a \\ 0 & -a & 0 \\ a & 0 & -a \end{pmatrix}. \quad \dots(32)$

Again m_0 is the contribution to the $0\nu\beta\beta$ decay.

For the mass matrix, it is also possible to give a simple parametrization which guarantees a maximal mixing for the atmospheric neutrinos and gives the $0\nu\beta\beta$ decay

$$\mathcal{M}_V = \begin{pmatrix} m_{ee} & a & a \\ a & b+c & b-c \\ a & b-c & b+c \end{pmatrix}. \quad \dots(33)$$

The $0\nu\beta\beta$ decay determines the element m_{ee} . The textures A1, A2 and A3 are the limiting cases with $a, c \ll b$; $a \ll b = c$ and $a \ll b = -c$, respectively. The CHOOZ constraint is satisfied and the solar neutrino solutions are obtained with suitable choice of a and c .

In the case of four generations the analysis becomes more involved. Only the mass matrices with minimum number of parameters have been studied in this context³⁸. A generalization of these results to include the $0\nu\beta\beta$ decay contribution can be written in the basis $[v_e, v_\mu, v_\tau, v_s]$ as³⁷

$$\mathcal{M}_{4\nu} = \begin{pmatrix} m & 0 & a & d \\ 0 & c & b & 0 \\ a & b & 0 & 0 \\ d & 0 & 0 & -m \end{pmatrix}. \quad \dots(34)$$

We can further economize by identifying two parameters $m = d$, making it effectively a four-parameter mass matrix. Here m determines the $0\nu\beta\beta$ contribution. The oscillation between the states v_e and v_s explains the solar neutrino problem. The mixing angle now becomes, $\sin^2 2\theta_{sol} = d^2/(m^2 + d^2)$. A simple choice of $d = m$ gives $\sin^2 2\theta_{sol} = 0.5$, which is consistent with present data. Restricting ourselves to $c \ll b$ ensures a maximal mixing between v_μ and v_τ , as required by the atmospheric neutrinos. The mass squared difference for the atmospheric neutrinos is given by $2bc$ and that for the LSND is $b^2 - d^2 - m^2$. There are no simple expressions for the mass squared difference required for the solar neutrinos and the mixing angle for the LSND result. Numerically these predictions come out as required. If LSND result is not confirmed by the MiniBOONE experiment, then this solution will become redundant considering the fact that SNO and KamLAND disfavors a sterile neutrino.

7 Models of Neutrino Masses

The $0\nu\beta\beta$ decay implies that neutrinos are Majorana particles. If neutrinos are Majorana particles, then the

smallness of the neutrino mass can also be explained naturally in any extension of the standard model. Consider the most general dimension-5 effective operator which can contribute to the Majorana masses of the neutrinos³⁹

$$\mathcal{L}_{Maj} = \Lambda^{-1} (v\phi^\circ - e\phi^+)^2, \quad \dots(35)$$

where Λ is the lepton number violating heavy scale in the theory and ϕ is the usual Higgs doublet scalar, transforming as $[1, 2, 1/2]$ under $SU(3)_c \times SU(2)_L \times U(1)_Y$. After the electroweak symmetry breaking this term will give Majorana masses to the neutrinos. A moderately high lepton number breaking scale can then give a small Majorana neutrino mass naturally. There are several mechanisms where this idea is realized in different ways³⁹. We shall demonstrate only a couple of simple realizations.

In the standard model neutrinos are massless. In the see-saw mechanism of the neutrino masses, one introduces a right-handed neutrino for each generation⁴⁰. So, corresponding to the three left-handed neutrinos, there are now three right-handed neutrinos $N_{iR}, i = 1, 2, 3$. The mass term for the neutrinos is then

$$\begin{aligned} \mathcal{L}_{mass} &= m_{D\alpha i} v_\alpha N_i^c + M_i N_i^c N_i^c \\ &= (v_\alpha \quad N_i^c) \begin{pmatrix} 0 & m_{D\alpha i} \\ m_{D\alpha i} & M_i \end{pmatrix} \begin{pmatrix} v_\alpha \\ N_i^c \end{pmatrix} \end{aligned} \quad \dots(36)$$

Here m_D comes from the standard model Higgs vacuum expectation values, so it is of the order of charged lepton masses. But the Majorana mass is very large, $M_M \sim 10^{10}$ GeV. Thus assuming $m_D \ll M$, the eigenvalues of this mass matrix then become,

$$m_1 = -\frac{m_D^2}{M} \quad \text{and} \quad m_2 = M. \quad \dots(37)$$

We then get a light neutrino with mass ~ 1 eV, which is mostly the left-handed neutrino with a small mixing $\tan \theta = \frac{2M}{m_D}$ with the right-handed neutrino. For the neutrinoless double beta decay there will be contributions from both the states with masses m_1 and m_2 . Since the effective mass is the electron neutrino mass term, now $\langle m \rangle = m_{ee} = \sum_{i=1}^2 U_{ei}^2 m_i F(m_i, A)$. Since m_2 is very large, $F(m_2, A) = 0$ and we get $\langle m \rangle \sim m_1$. Thus, for studying the $0\nu\beta\beta$ decay we may work with the effective mass matrix of the light neutrino states only.

We shall now consider a different realization of the effective operator (35), where the standard model is extended to include a triplet Higgs scalar^{41, 42}. For this triplet ξ to couple to two left-handed leptons, it should transform as symmetric product $\phi\phi$, which fixes the hypercharge of the triplet to be $Y = 1$. Thus under $SU(3)_c \times SU(2)_L \times U(1)_Y$ it transforms as $[1, 3, +1]$. The three components of the triplet are $\xi \equiv [\xi^{++}, \xi^+, \xi^0]$. Its couplings to the leptons and to the usual standard model Higgs doublets ϕ break lepton number,

$$\mathcal{L}_{Yuk} = f_{ij} \xi \ell_i \ell_j + \mu \xi^\dagger \phi \phi. \quad \dots(38)$$

We consider⁴¹ the possibility $\mu \neq 0$ ($\mu = 0$ models⁴² are ruled out by LEP data). The neutral component of ξ will acquire a vev during the $SU(2)_L$ symmetry breaking $u = \frac{-\mu v^2}{M^2}$, where M is the mass of the triplet Higgs ξ . The lepton number is broken explicitly at a very high scale and $M \sim \mu \sim 10^{13}$ GeV. So, there are no Goldstone bosons corresponding to the broken lepton number symmetry. The mass of the left-handed neutrinos are then given by

$$m_{ij}^v = f_{ij} u = -f_{ij} \frac{\mu v^2}{M^2}, \quad \dots(39)$$

which is of the order of \sim eV. The neutrino mass matrix is now directly proportional to the Yukawa couplings f_{ij} . In particular, the (11) element determines the amount of the neutrinoless double beta decay.

8 Leptogenesis

The neutrinoless double beta decay implies that neutrinos are Majorana particles and lepton number is violated at some scale Λ . It has been pointed out that this lepton number violation can solve the problem of the baryon asymmetry of the universe^{43, 44}, *i.e.*, why there is more matter in the universe compared to the anti-matter. To establish this connection between the neutrinoless double beta decay and baryogenesis, we shall explain here how baryogenesis could be explained in the two models of neutrino masses discussed in the previous section.

Except for the neutrinoless double beta decay, all other experiments with neutrinos could be explained with Dirac neutrinos and without any lepton number violation. Thus the neutrinoless double beta decay implies that there is lepton number violation and

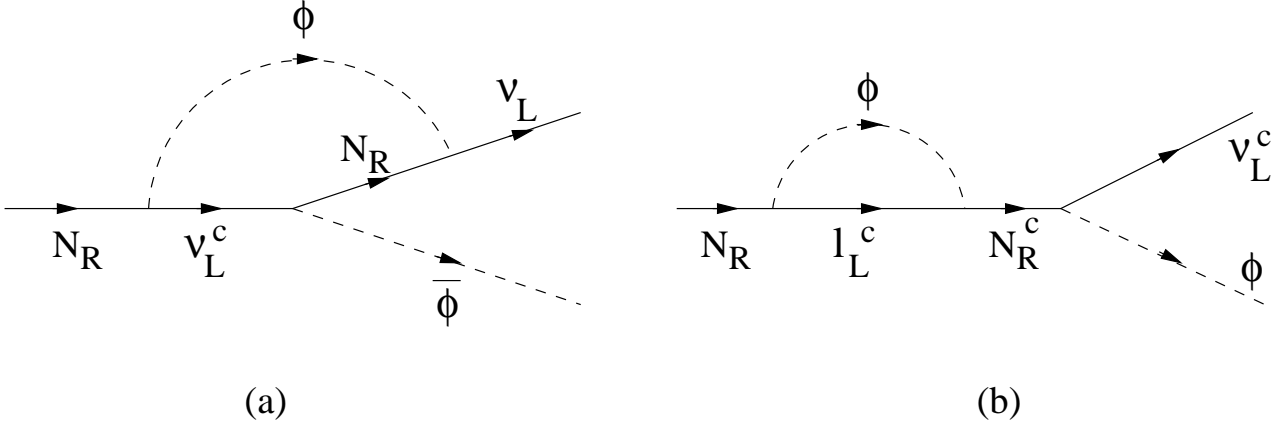


Fig. 9 One loop (a) vertex and (b) self energy diagrams, which interferes with the tree level right-handed neutrino decays to produce CP violation

hence $(B - L)$ violation in nature at some large scale $\Lambda \sim 10^{10}$ GeV. There are $SU(2)_L$ global anomalies in the standard model, which breaks both B and L by equal amount, leaving $B - L$ unbroken. At finite temperature, during the period $10^2 < T < 10^{12}$ GeV, the $(B + L)$ violating interactions becomes too strong in the presence of some static topological field configuration called the sphalerons⁴⁵. The $(B + L)$ asymmetry of the universe is then washed out by the sphalerons before the electroweak phase transition. However, the $(B - L)$ asymmetry of the universe required by the neutrinoless double beta decay would survive till low energy and the present day baryon asymmetry of the universe will be related to this $(B - L)$ asymmetry of the universe after the electroweak phase transition^{43, 44}.

The baryon asymmetry of the universe generated in this process requires lepton number violating processes at some large scale. These processes should also violate C and CP and should take place away from equilibrium. All these three conditions are satisfied in the models of neutrino masses we discussed in the previous section.

In the see-saw mechanism of neutrino masses, lepton number violation comes from the Majorana masses of the heavy right-handed neutrinos. The right-handed neutrinos can now decay into a lepton ($N_{Ri} \rightarrow \ell_{jL} + \bar{\phi}$) and also an antilepton ($N_{Ri} \rightarrow \ell_{jL}^c + \phi$) violating lepton number. The \mathcal{CP} violation comes from an interference of the tree level decays of the right-handed neutrinos and the one loop diagrams involving right-handed neutrinos of another generation. There are two sources of CP violation:

(i) vertex diagram^{43, 46} of Fig.9a, which interferes with the tree level decays of the right-handed

neutrinos $N_R \rightarrow v_L + \bar{\phi}$. This is similar to the CP violation coming from the penguin diagram in K -decays.

(ii) self energy diagram⁴⁷ of Fig.9b, which interferes with the tree level decays of the right-handed neutrinos $N_R \rightarrow v_L^c + \phi$ to produce CP violation. This is similar to the CP violation in $K - \bar{K}$ oscillation, entering in the mass matrix of the heavy Majorana neutrinos.

The vertex diagrams are the ones which are usually considered for the asymmetry in decays of particles, which is given by,

$$\begin{aligned} \delta &= \frac{\Gamma(N \rightarrow \ell \phi^\dagger) - \Gamma(N \rightarrow \ell^c \phi)}{\Gamma(N \rightarrow \ell \phi^\dagger) + \Gamma(N \rightarrow \ell^c \phi)} \\ &= -\frac{1}{8\pi} \frac{M_1 M_2}{M_2^2 - M_1^2} \frac{\text{Im}[\sum_\alpha (h_{\alpha 1}^* h_{\alpha 2}) \sum_\beta (h_{\beta 1}^* h_{\beta 2})]}{\sum_\alpha |h_{\alpha 1}|^2} \end{aligned} \quad \dots (40)$$

In this expression it has been assumed that the main contribution to the asymmetry comes from the lightest right-handed neutrino (N_1) decay. When the mass difference is large compared to the width, the CP asymmetry generated through the mixing of the heavy neutrinos is the same as the vertex correction as given by eq.40. These two contributions then add up to produce the final lepton asymmetry of the universe.

The asymmetry is generated when the lightest right-handed neutrino (say N_1) decay, and it satisfy the out-of-equilibrium condition, which is,

$$\frac{|h_{\alpha 1}|^2}{16\pi} M_1 < 1.7 \sqrt{g^*} \frac{T^2}{M_p} \quad \text{at } T = M_1 \dots (41)$$

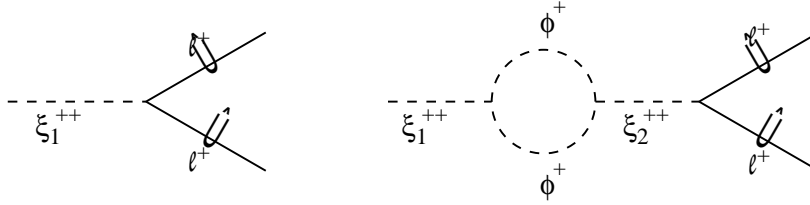


Fig. 10 The decay of $\xi_1^{++} \rightarrow l^+l^+$ at tree level and in one-loop order, whose interference gives \mathcal{CP} violation.

which gives a bound on the mass of the lightest right-handed neutrino to be $m_{N_1} > 10^7 \text{ GeV}^{48}$. The lepton asymmetry thus generated is same as the $(B-L)$ asymmetry of the universe, since there is no primordial baryon asymmetry at this time. The sphaleron +interactions now convert this $(B-L)$ asymmetry to a baryon asymmetry of the universe.

In the triplet Higgs mechanism of neutrino masses⁴¹, the decays of the heavy triplet Higgs scalars ($\xi_a, a = 1, 2$, two scalars are required for \mathcal{CP} violation) violate lepton number

$$\xi_a^{++} \rightarrow \begin{cases} l_i^+ l_j^+ & (L = -2) \\ \phi^+ \phi^+ & (L = 0) \end{cases} \dots (40)$$

These interactions can generate a lepton asymmetry of the universe if they are slow enough and there is required \mathcal{CP} violation. In this case there are no vertex corrections which can introduce CP violation. The only source of CP violation is the self energy diagrams of Fig.10, when there are more than one triplet Higgs.

In the presence of the one loop diagram, the mass matrices M_a^2 and M_a^{*2} become different. This implies that the rate of $\xi_b \rightarrow \xi_a$ no longer remains to be the same as $\xi_b^* \rightarrow \xi_a^*$. Since by CPT theorem $\xi_b^* \rightarrow \xi_a^* \equiv \xi_a \rightarrow \xi_b$, it means

$$\Gamma[\xi_a \rightarrow \xi_b] \neq \Gamma[\xi_b \rightarrow \xi_a]. \dots (41)$$

This is a different kind of CP violation compared to the CP violation in models with right-handed neutrinos or all other models of decays of scalars, where the vertex diagrams are mainly considered for the generation of a baryon asymmetry of the universe. This \mathcal{CP} violation is analogous to the lepton number conserving \mathcal{CP} violation, which may enter in the neutrino oscillation experiments.

If we consider that the ξ_2 is heavier than ξ_1 , then the decay of ξ_1 will generate an lepton asymmetry

given by,

$$\begin{aligned} \delta &= \frac{\Gamma(\xi \rightarrow \ell\ell) - \Gamma(\xi^\dagger \rightarrow \ell^c \ell^c)}{\Gamma(\xi \rightarrow \ell\ell) + \Gamma(\xi^\dagger \rightarrow \ell^c \ell^c)} \\ &= \frac{\text{Im} \left[\mu_1 \mu_2^* \sum_{k,l} f_{1kl} f_{2kl}^* \right]}{8\pi^2 (M_1^2 - M_2^2)} \left[\frac{M_1}{\Gamma_1} \right] \dots (42) \end{aligned}$$

In this model the out-of-equilibrium condition is satisfied when the masses of the triplet Higgs scalars are heavier than 10^{13} GeV . The lepton asymmetry thus generated after the Higgs triplets decayed away would then generate a baryon asymmetry of the universe $n_B/s \sim 10^{-10}$, as desired.

9 Summary

We discussed the implications of the neutrinoless double beta decay experiment¹⁰. The existing upper bound on the neutrino mass already gave us several interesting bounds on all the lepton number violating processes. There are constraints on the masses and couplings of the heavy right-handed neutrinos, lept-quarks, and several other scalar bilinears, which are much stronger than the direct constraints coming from the accelerators. In supersymmetric models also there are strong constraints coming from the neutrinoless double beta decay. All lepton number violating and R-parity violating couplings are also strongly constrained. The reported evidence for the neutrinoless double beta decay implies lepton number violation in nature and that the neutrinos are Majorana particles. This will give us natural scenarios of generating small neutrino masses and allow us to explain the baryon asymmetry of the universe. It further constrains the neutrino mass matrices very severely and rule out several classes of models.

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