Particle Physics Implications of Neutrinoless Double Beta Decay∗
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1. Introduction

In the standard electroweak model of Glashow, Weinberg and Salam, the absence of the righthanded neutrinos and the existence of an exact accidental global $B-L$ symmetry guarantees that the neutrinos are massless to all orders in perturbation theory. Any experimental evidence for a non-zero neutrino mass therefore constitutes evidence for new physics beyond the standard model and will be a major step towards a deeper understanding of new forces in nature[1]. Among the many experiments that are under way at this moment searching directly or indirectly (e.g. via neutrino oscillations) for neutrino masses, one of the most important ones is the search for neutrinoless double beta decay. This process is allowed only if the neutrino happens to be its own antiparticle ( Majorana neutrino) as is implied by many extensions of the standard model. In fact there is a well-known theorem[2] that states that any evidence for neutrinoless double beta decay is an evidence for nonzero Majorana mass for the neutrinos. It is of course a much more versatile probe of new physics as we will discuss in this article. The point is that since $\beta\beta_{0\nu}$ decay changes lepton number ($L_e$) by two units any theory that contains interactions that violate electron lepton number $L_e$ can in principle lead this process. This therefore reflects the tremendous versatility of $\beta\beta_{0\nu}$ decay as a probe of all kinds of new physics beyond the standard model. Indeed we will see that already very stringent constraints on new physics scenarios such as the left-right symmetric models with the see-saw mechanism[3] and supersymmetric models with R-parity violation[4], scales of possible compositeness of leptons etc are implied by the existing experimental limits[5] on this process. For a more detailed discussion of the theoretical situation than is possible here, see[6]. For an update of the experimental situation, both ongoing and in planning stage, see[7].

This talk is organized as follows: In section 2, I discuss the basic mechanisms for neutrinoless double beta decay ; in section 3, the implications of the present limits on the lifetime for neutrinoless double beta decay for neutrino mixings are discussed; in part section 4, I go on to discuss the kind of new physics scenarios that can be probed by $\beta\beta_{0\nu}$ decay and the constraints on the parameters of the new physics scenarios implied by present data.

2. Mechanisms for $\beta\beta_{0\nu}$ decay

As is wellknown, if the neutrino is its own antiparticle, the conventional four-Fermi interaction can lead to neutrinoless double beta decay via the diagram in Fig. 1. In physics scenarios beyond the standard models, if there are heavy Majorana fermions interacting with the electrons, diagrams

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similar to Fig. 1 with neutrino line replaced by the Majorana fermions can also lead to $\beta\beta$ decay. Examples of such particles abound in literature: right-handed neutrino, photino, gluino to mention a few popular ones.

![Feynman diagram](image)

Figure 1. Feynman diagram involving neutrino majorana mass that contributes to $\beta\beta$ decay.

One could therefore give an arbitrary classification of the mechanisms for $\beta\beta$ decay into two kinds: (A) one class that involves the exchange of light neutrinos; and (B) the second class that involves heavy fermions or bosons. Furthermore, there are two distinct mechanisms for light neutrino exchange contributions: (a) helicity flip light neutrino mass mechanism and (b) helicity nonflip vector-vector or vector-scalar mechanism. In case (a), one can write the amplitude $A_{\beta\beta}$ for neutrinoless double beta decay to be:

$$A_{\beta\beta}^{(m)} \simeq \frac{G_F^2}{2} \langle m_\nu | \frac{\eta}{\mathbf{k} \cdot \mathbf{r}} | N_{ucl}\rangle$$  \hspace{1cm} (1)

whereas in case (b), it looks like:

$$A_{\beta\beta} \simeq \frac{G_F^2}{2} \langle \eta | \mathbf{k} \cdot \mathbf{r} | N_{ucl}\rangle$$  \hspace{1cm} (2)

To extract neutrino mass implications for neutrinoless double beta decay, we need to note the explicit form of $< m_\nu >$:

$$< m_\nu > = \Sigma_i U_{ei}^2 m_i$$  \hspace{1cm} (3)

where $U_{ei}$ are the mixing matrix elements for the electron neutrino with the other neutrinos. Therefore a constraint on the $< m_\nu >$ can be converted into constraints on the neutrino mixings involving the first generation. Incidentally, one can also write $< m_\nu > = m_{ee}$ where $m_{ee}$ is the $ee$ entry of the neutrino mass matrix in the weak basis. Thus any theory which has zero entry in the $ee$ location leads to vanishing neutrinoless double beta decay even if the neutrino is a Majorana particle.

It is important to remark that these kind of light neutrino exchange diagrams always lead to a long range neutrino potential inside the nucleons and therefore, crudely speaking the two nucleons ”far” from each other can contribute in the double beta decay. This has important implications for the evaluation of the nuclear matrix element, an important subject we do not discuss here. We will instead use an effective momentum to parameterize the effect of the nuclear matrix element calculations (we will roughly choose $p_{\text{eff}} \approx 50$ MeV). The width for double beta decay amplitude is given by

$$\Gamma_{\beta\beta} \simeq \frac{Q^5 |A|^2}{60\pi^3}$$  \hspace{1cm} (4)

Here, $Q$ is the available energy for the two electrons. Using the present most stringent limit on $\tau_{\beta\beta} \geq 1.1 \times 10^{25}$ years obtained for $^{76}\text{Ge}$ by the Heidelberg-Moscow group, one can obtain the upper limit on the width to be $\Gamma_{\beta\beta} \leq 3.477 \times 10^{-57}$ GeV; using Eq. (4), $A$ for the light neutrino contribution, $Q \simeq 2$ MeV and $p_F \simeq 50$ MeV, one gets a rough upper limit of 0.7 eV for the neutrino mass. A more careful estimate leads to

$$< m_\nu > \leq 0.46 \text{ eV} \hspace{1cm} \eta \leq 10^{-8}$$  \hspace{1cm} (5)
The second class of mechanisms consists of exchange of heavy particles which often arise in physics scenarios beyond the standard model. In the low energy limit, the effective Hamiltonian that leads to $\beta\beta$ decay in these cases requires point interaction between nucleons; as a result, in general the nuclear matrix elements are expected to be smaller due to hard core repulsive nuclear potential; nevertheless, a lot of extremely useful information have been extracted about new physics where these mechanisms operate. Symbolically, such contributions can arise from effective Hamiltonians of the following type (we have suppressed all gamma matrices as well as color indices):

$$H^{(1)} = G_{\text{eff}} \bar{\pi} \Gamma d \Gamma F + h.c.$$  \hspace{1cm} (6)

or

$$H^{(2)} = \lambda_\Delta \left( \frac{1}{M^4} \bar{\pi} \Gamma d \Gamma d + e^- e^- \right) \Delta^{++} + h.c.$$  \hspace{1cm} (7)

Here $F$ represents a neutral majorana fermion such as the right-handed neutrino ($N$) or gluino $\tilde{G}$ or photino $\tilde{\gamma}$ and $\Delta^{++}$ represents a doubly charged scalar or vector particle. In the above equations, the coupling $G_{\text{eff}}$ has dimension of $M^{-2}$ and $\lambda_\Delta$ is dimensionless. The possibility of the doubly charged scalar contribution to $\beta\beta$ was first noted in [11] and have been discussed subsequently in [12]. The contributions to neutrinoless double beta decay due to the above interactions lead to $\beta\beta$ amplitudes of the form:

$$A_{\beta\beta}^{(F)} \simeq G_{\text{eff}}^2 \frac{1}{M_F} (p_{\text{eff}})^3$$  \hspace{1cm} (8)

and

$$A_{\beta\beta}^{\Delta} \simeq \left( \frac{\lambda_\Delta^2}{M^4 M_\Delta^2} \right) (p_{\text{eff}})^3$$  \hspace{1cm} (9)

Here again we have crudely replaced all nuclear effects by the effective momentum parameter $p_{\text{eff}}$. If we choose $p_{\text{eff}} \simeq 50$ MeV, then the present lower limit on the lifetime for $^{76}\text{Ge}$ decay leads to a crude upper limit on the effective couplings as follows:

$$G_{\text{eff}} \leq 10^{-7} \left( \frac{M_F}{100 \text{ GeV}} \right)^{\frac{3}{2}}$$  \hspace{1cm} (10)

and

$$\lambda_\Delta \leq 10^{-3} \left( \frac{M}{100 \text{ GeV}} \right)^{\frac{1}{2}}$$  \hspace{1cm} (11)

In the second equation above, we have set $M = M_\Delta$. Note that these limits are rather stringent and therefore have the potential to provide useful constraints on the new physics scenarios that lead to such particles.

3. Implications for neutrino masses and mixings

This conference watched the history of neutrino physics take a remarkable new turn. Convincing evidence was presented by the Super-Kamiokande collaboration for the existence of neutrino oscillation of the atmospheric muon neutrinos to either $\nu_\tau$ or a sterile neutrino. Using data both in the sub-GeV and multi-GeV energy range for the electron and the muon neutrinos as well as the zenith angle dependence of the muon data, the present fits at 90% confidence level seem to imply the following values for the oscillation parameters $\Delta m^2$ and $\sin^2 2\theta$: $4 \times 10^{-4} \leq \Delta m^2 \leq 5 \times 10^{-3}$ eV$^2$ with $\sin^2 2\theta$ between .8 to 1 [13]. The possibility of $\nu_\mu - \nu_e$ oscillation as an explanation of the atmospheric anomaly seems to run into conflict with the recent CHOOZ [13] experiments. Neutrino oscillation also seems to be the only way to understand the deficit of the solar neutrinos [14]. The detailed oscillation mechanism in this case is however is unclear. The three possibilities are: a) Small-angle MSW [15], $\Delta m^2_{\text{sol}} \approx 6 \times 10^{-6}$ eV$^2$, $\sin^2 2\theta_{\text{sol}} \approx 7 \times 10^{-3}$; b) Large-angle MSW, $\Delta m^2_{\text{atm}} \approx 9 \times 10^{-6}$ eV$^2$, $\sin^2 2\theta_{\text{atm}} \approx 0.6$; c) Vacuum oscillation, $\Delta m^2_{\text{vac}} \approx 10^{-10}$ eV$^2$, $\sin^2 2\theta_{\text{vac}} \approx 0.9$. The data on neutrino energy distribution presented at this conference indicates a preference towards vacuum oscillation rather than MSW mechanism. Turning to the laboratory experiments, the LSND [16] collaboration has presented evidence in favor of a possible oscillation of $\overline{\nu}_\mu \rightarrow \overline{\nu}_e$ as well as $\nu_\mu - \nu_e$. The preferred $\Delta m^2$ range seems to be $24 \leq \Delta m^2_{\nu_{\mu-e}} \leq 10$ eV$^2$ with a mixing angle in the few percent range. As already mentioned, $\beta\beta$ gives only an upper bound of $< m_{\nu_e} > < .46$ eV.
Another effect of neutrino mass is in the arena of cosmology, where it not only effects whether the universe keeps expanding for ever or it eventually collapses onto itself, but it also determines the detailed manner in which structure formed in the early universe. This subject is in a constant state of flux due to new cosmological data coming in at a very rapid rate. But the idea that the present structure data may need a neutrino mass contribution to the dark matter is very much alive (see for instance Ref.[17] which seems to suggest that a total neutrino mass of 4-5 eV which contributes about 20% of the dark matter along with 70% cold dark matter and 10% baryon gives the best fit to the galaxy power spectrum data. This taken seriously would mean that $\Sigma m_{\nu_i} = 4 - 5$ eV).

With the above input information, if we stay within the minimal three neutrino picture, then the solar neutrino puzzle can be resolved by $\nu_e \rightarrow \nu_\mu$ oscillations and the atmospheric neutrino deficit by $\nu_\mu \rightarrow \nu_\tau$ oscillations and the LSND results cannot be accomodated. Note that these observables are controlled only by the mass square difference; on the other hand, the required hot dark matter implies that at least one or more of the neutrinos must have mass in the few eV range. It was pointed out[18] in 1993 that, in the minimal picture, this leads to a scenario, where all three neutrinos are nearly degenerate, with $m_{\nu_\mu} \approx 1.6$ eV. It is then clear that, in this case, in general there will be an observable amplitude for neutrinoless double beta decay mediated by the neutrino mass. In fact, if the limit on $\langle m_{\nu_\mu} \rangle$ is taken to be less than .47 eV as is implied for a certain choice of the nuclear matrix element, then the mixing must satisfy the constraint:

$$\Sigma_i U_{ei}^2 \simeq 0$$

(12)

Since each of the elements in the above sum is complex, the $U_{ei}^2$ form the three sides of a triangle[19]. Then using the unitarity relation for the $U$ matrix, it is clear that one must have $|U_{ei}| \leq 1/2$. On the other hand, the CHOOZ data for a general three neutrino oscillation picture implies that $4|U_{e3}|^2 (1 - |U_{e3}|^2) \leq .2$. These two constraints then imply that $|U_{e3}| \leq .2$. This is indeed an interesting constraint and rules out (provided of course $\Delta m_{32}^2 \geq 10^{-3}$ eV$^2$) a maximal mixing scenario for degenerate neutrinos that was proposed to reconcile sub-eV double beta decay neutrino mass limit i.e.[20]

$$U = \frac{1}{\sqrt{3}} \begin{pmatrix} 1 & \omega & \omega^2 \\ 1 & \omega^2 & \omega \\ 1 & 1 & 1 \end{pmatrix}$$

(13)

There is however another mixing pattern for the degenerate neutrino scenario which is consistent with both the CHOOZ experiment and the neutrinoless double beta decay bounds:

$$U = \begin{pmatrix} 1/\sqrt{2} & i/\sqrt{2} & 0 \\ 1/\sqrt{6} & -i/\sqrt{6} & -2/\sqrt{6} \\ 1/\sqrt{3} & -i/\sqrt{3} & 1/\sqrt{3} \end{pmatrix}$$

(14)

Other more general constraints for this case have been studied in several recent papers[21].

If we do not include the hot dark matter constraint, then there is no need to require that the neutrinos are degenerate in mass and one can live perfectly happily with a hierarchical pattern of neutrino masses as dictated by the simple type I seesaw formula. In that case, one can combine the atmospheric oscillation fits and the CHOOZ data to set an upper limit on $\langle m_{\nu_\mu} \rangle$ equal to $\sqrt{\Delta m_{\text{ATMOS}}^2 \sin^2 \theta_{\mu\tau}} \simeq .02$ eV[22]. Thus evidence for $\langle m_{\nu_\mu} \rangle$ above this value would be an indication that either the neutrino mass pattern is not hierarchical or that the atmospheric neutrino puzzle involves transition between $\nu_\mu$ and a sterile neutrino. Both of these are extremely valuable conclusions. The GENIUS proposal of the Heidelberg group[23] is expected to push the double beta decay limit to this level and could therefore test this conclusion.

4. Implications for physics beyond the standard model:

Let us now discuss the constraints implied by neutrinoless double beta decay searches on the new physics scenarios beyond the standard model. Let us first consider the the neutrino mass mechanism. Any theory which gives the electron neutrino a significant ($\simeq$ eV) Majorana mass or any other species (e.g. $\nu_\mu$ or $\nu_\tau$) a large enough mass and mixing angle with the $\nu_e$ so
that $U_{e2}^2 m_{\nu_e}$ is of order of an electron volt will make itself open to testability by the $\beta\beta_{0\nu}$ decay experiment. There are many theories with such expectations for neutrinos. Below I described two examples: (i) the singlet majoron model and (ii) the left-right symmetric model. Both these models are intimately connected with ways to understand the small neutrino mass in gauge theories.

### 4.1. The singlet majoron model:

This model\cite{24} is the simplest extension of the standard model that provides a naturally small mass for the neutrinos by employing the the see-saw mechanism\cite{25}. It extends the standard model by the addition of three right-handed neutrinos and the addition of a single complex Higgs field $\Delta$ which is an $SU(2)_L \times U(1)_Y$ singlet but with a lepton number +2. There is now a Dirac mass for the neutrinos and a Majorana mass for the right handed neutrinos proportional to the vacuum expectation value (vev) $\langle \Delta \rangle \equiv \nu_R$. This leads to a mass matrix for the neutrinos with the usual see-saw form:

$$M = \begin{pmatrix} 0 & m_D \cr m_D^T & f_{\nu R} \end{pmatrix}$$ (15)

This leads to both the light and heavy (right-handed) neutrinos being Majorana particles with the mutual mass relation being given by the see-saw formula:

$$m_{\nu_i} \simeq m_{1D}(M_{1R}^{-1})m_D^T$$ (16)

where we have ignored all mixings and $M_{1R} \simeq f_{\nu R}$ denote the masses of the heavy right-handed neutrinos. It is clear that the electron neutrino mass can be in the electron-volt range if the values of $m_{1D}$ are chosen to be of similar order of magnitude to the electron mass. In fact, for $m_{1D} = m_e$, and $m_{1R} = 250$ GeV, one gets $m_{\nu_e} = 1$ eV which is the range of masses being probed by the ongoing and proposed $\beta\beta_{0\nu}$ experiments.

More importantly, this class of models leads to the new neutrinoless double beta decay process with majoron emission\cite{28} which has a very different electron energy distribution than either 0$\nu$ or 2$\nu$ double beta decays. The relevant Feynman diagram is same that in Fig. 1 with a majoron line emanating from the light neutrino in the middle. The majoron coupling $g_{\nu \nu X}$ then replaces the neutrino mass in the $\beta\beta_{0\nu}$ amplitude. This observation has led to a considerable amount of experimental effort into searching for the majoron emitting double beta decay and limits at the level of $g_{\nu \nu X} \leq 10^{-5}$ are presently available.

A relevant question is whether majoron couplings at the level measurable are expected in reasonable extensions of the standard model. There have been extensive studies of this question and is beyond the scope of this review. But it is of interest to note that in the simplest singlet majoron model, one expects $g_{\nu \nu X} \simeq \Sigma_a m_{\nu a}^2 M_{a}^{-2} g_{\alpha a}$. In the absence of any mixings, this is proportional to $m_{\nu_a}/M_{\nu_1}$ which is expected to be of order $10^{-11}$ for an eV $\nu_e$ and 100 GeV for the $B-L$ breaking scale. However, if the $\nu_e$ mass is in the MeV range as is allowed by LEP analysis, this coupling could easily be in the $10^{-5}$ to $10^{-6}$ range which is clearly in the range accessible to experiments.

### 4.2. Left-right symmetric models:

Let us now consider the minimal left-right symmetric model with a see-saw mechanism for neutrino masses as described in\cite{3}. Below, we provide a brief description of the structure of the model. The three generations of quark and lepton fields are denoted by $Q_a^T \equiv (u_a, d_a)$ and $\Psi_a^T \equiv (\nu_a, e_a)$ respectively, where $a = 1, 2, 3$ is the generation index. Under the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, they are assumed to transform as $\Psi_{a, L} \equiv (1/2, 0, -1)$ and $\Psi_{a, R} \equiv (0, 1/2, -1)$ and similarly for the quarks denoted by $Q^T \equiv (u, d)$. In this model, there is a right-handed counterpart to the $W^\pm$ to be denoted by $W^+_R$. Their gauge interactions then lead to the following expanded structure for the charged weak currents in the model for one generation prior to symmetry breaking (for our discussion, the quark mixings and the higher generations are not very important; so we will ignore them in what follows.)

$$L_{\text{wk}} = \frac{g}{2\sqrt{2}} [W_{\mu L} J^\mu_L + L \rightarrow R]$$ (17)

where $J^\mu_L = (\overline{d} \gamma^\mu (1 - \gamma_5) u + \overline{e} \gamma^\mu (1 - \gamma_5) \nu_e)$
The Higgs sector of the model consists of the bi-doublet field \( \phi \equiv (1/2, 1/2, 0) \) and triplet Higgs fields: \( \Delta_L(1, 0, +2) \oplus \Delta_R(0, 1, +2) \).

The Yukawa couplings for the lepton sector which are invariant under gauge and parity symmetry can be written as:

\[
\mathcal{L}_Y = \overline{\Psi}_L h^{\ell} \phi \Psi_R + \overline{\Psi}_L h^{\ell} \phi \Psi_R + \Psi_L^T \gamma^\tau \tilde{\Delta}_L C^{-1} \Psi_L + \text{h.c.} \tag{18}
\]

where \( h, \tilde{h} \) are hermitian matrices while \( f \) is a symmetric matrix in the generation space. \( \Psi \) and \( \xi \) are neutrino with a small strength proportional to the charged current Lagrangian, we see that the right-handed neutrinos whereas \( < \Delta^0_R > < \Delta^0_L > 0 \)).

The gauge symmetry is spontaneously broken by the vacuum expectation values:

\[
< \Delta^0_R > = v_R \quad \text{and} \quad < \phi > = \left( \begin{array}{cc}
\kappa & 0 \\
0 & \kappa' \\
\end{array} \right).
\]

As usual, \( < \phi > \) gives masses to the charged fermions and Dirac masses to the neutrinos whereas \( < \Delta^0_R > \) leads to the seesaw mechanism for the neutrinos in the standard way. For one generation the seesaw matrix is in the form \( m_\nu \approx m_f^2/fv_R \) and leads as before to a light and a heavy state as discussed in the previous section. For our discussion here it is important to know the structure of the light and the heavy neutrino eigenstates:

\[
\nu \equiv \nu_e + \xi N_e \\
N \equiv N_e - \xi \nu_e \tag{19}
\]

where \( \xi \approx \sqrt{m_\nu/m_N} \) and is therefore a small number. Substituting these eigenstates into the charged current Lagrangian, we see that the right-handed \( W_R \) interaction involves also the light neutrino with a small strength proportional to \( \xi \).

To second order in the gauge coupling \( g \), the effective weak interaction Hamiltonian involving both the light and the heavy neutrino becomes:

\[
H_{\text{ewk}} = \frac{G_F}{\sqrt{2}} (\overline{\nu} \gamma^\mu (1 - \gamma_5) \nu)d[\overline{\nu} \gamma_\mu(1 - \gamma_5)] + \xi(\frac{m_\nu^2}{m_{W_R}})(1 + \gamma_5)\nu + \xi(1 + \gamma_5)N
\]

\[
+ \frac{G_F}{\sqrt{2}} \left( \frac{m_\nu^{24}}{m_{W_R}} \right) (\overline{\nu} \gamma^\mu (1 + \gamma_5) \nu \overline{\nu} \gamma_\mu(1 + \gamma_5)N) + \text{h.c.} \tag{20}
\]

From Eq. (20), we see that there are several contributions to the \( \beta\beta_{0v} \) process. Aside from the usual neutrino mass diagram (Fig.1), there is a contribution due to the wrong helicity admixture with \( 0 < \xi \left( \frac{m_\nu^{24}}{m_{W_R}} \right) \) and there are contributions arising from the exchange of heavy right-handed neutrinos. This last contribution is given by:

\[
A_{\beta\beta}^{(R)} \approx \frac{G_F^2}{2} \left( \frac{m_{W_L}^4}{m_{W_R}^2} + \xi^2 \right) \frac{1}{m_N} \tag{21}
\]

The present limits on neutrinoless double beta decay lifetime then imposes a correlated constraint on the parameters \( m_{W_R} \) and \( m_N \). If we combine the theoretical constraints of vacuum stability then, the present \( ^{76}\text{Ge} \) data provides a limit on the masses of the right handed neutrino (\( N_e \)) and the \( W_R \) of 1 TeV, which is a rather stringent constraint. We have of course assumed that the leptonic mixing angles are small so that there is no cancellation between the parameters.

Finally, the Higgs sector of the theory generates two types of contributions to the \( \beta\beta_{0v} \) decay. One arises from the coupling of the doubly charged Higgs boson to electrons (see Fig.2). The amplitude for the decay is same as in Eq. (6) except we have \( \lambda_\Delta = f_{11} \) and

\[
\lambda_\Delta M^3 \approx 2^{7/4} G_F^{3/2} \left( \frac{m_{W_L}}{m_{W_R}} \right)^3 \tag{22}
\]

Using this expression, we find that the present \( ^{76}\text{Ge} \) data implies that (assuming \( m_{W_R} \geq 1 \text{ TeV} \))

\[
M_{\Delta^{++}} \geq \sqrt{f_{11} \times 80 \text{GeV}} \tag{23}
\]

A second type Higgs induced contribution arises from the mixing among the charged Higgs fields in \( \phi \) and \( \Delta_L \) which arise from the couplings in the Higgs potential, such as \( \text{Tr}(\Delta_L \phi \Delta_R^\dagger) \) after the full gauge symmetry is broken down to \( U(1)_{em} \). Let us denote this mixing term by an angle \( \theta \). This will contribute to the four-Fermi interaction of the form given by the \( \epsilon_1^{ee} \) term with

\[
\epsilon_1^{ee} \approx \frac{h_u f_{11} \sin 2\theta \times}{4 \sqrt{2} G_F M_{H^+}^2}, \tag{24}
\]

where we have assumed that \( H^+ \) is the lighter of the two Higgs fields. We get \( h_u f_{11} \sin 2\theta \leq 6 \times \)
\(10^{-9}(M_{H^+}/100 \, \text{GeV})^2\), which is quite a stringent constraint on the parameters of the theory. To appreciate this somewhat more, we point out that one expects \(h_u \approx m_u/m_W \approx 5 \times 10^{-5}\) in which case, we get an upper limit for the coupling of the Higgs triplets to leptons \(f_{11}\sin 2\theta \leq 10^{-4}\) (for \(m_{H^+} = 100 \, \text{GeV}\)). Taking a reasonable choice of \(\theta \sim M_W/\sqrt{m_W m_R} \sim 10^{-1}\) would correspond to a limit \(f_{11} \leq 10^{-3}\). Limits on this parameter from analysis of Bhabha scattering is only of order .2 or so for the same value of the Higgs mass.

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An interesting recent development is that once one supersymmetrizes the seesaw version of the left-right model just described, allowed values for the right handed scale get severely restricted by the requirement that the ground state of the theory conserve electric charge. There are only two allowed domains for \(M_{W_R}\): (i) if the ground state breaks R-parity, there is an upper limit on the \(W_R\) scale of about \(\leq 10 \, \text{TeV}\). Since in this case, R-parity is spontaneously broken R-parity violating interactions conserve baryon number and the theory therefore is much improved in the sense of naturalness over the MSSM. What is interesting is that the GENIUS experiment can then completely scan the allowed range of this model. On the other hand, if R-parity is conserved, there must be a lower limit on \(M_{W_R}\) of about \(10^{10} \, \text{GeV}\). In this case also there is a contribution to \(\beta\beta^0\nu\) decay coming from the light doubly charged Higgs boson in the same manner described above. This contribution scales like \(V_R^{-2}\) in the amplitude. Thus as the limits on neutrinoless double beta decay improve, at some point they will not only imply that the \(W_R\) mass is not only bigger that \(10^{10} \, \text{GeV}\) or so; but they can also continue to improve this lower limit due to the contribution from the doubly charged Higgs boson whose mass is directly proportional to the square of \(v_R\).

4.3. MSSM with R-parity violation:

The next class of theories we will consider is the supersymmetric standard model. As is well-known, the minimal supersymmetric standard model can have explicit violation of the R-symmetry (defined by \((-1)^{B+L+2S}\)), leading to lepton number violating interactions in the low energy Lagrangian. The three possible types of couplings in the superpotential are:

\[
W' = \lambda_{ijk} L_i L_j E_{c_k}^c + \lambda'_{ijk} L_i Q_j D_{c_k}^c + \lambda''_{ijk} U_{c_i}^c D_{c_j}^c D_{c_k}^c .
\]

Here \(L, Q\) stand for the lepton and quark doublet superfields, \(E_c^c\) for the lepton singlet superfield and \(U^c, D^c\) for the quark singlet superfields. \(i, j, k\) are the generation indices and we have \(\lambda_{ijk} = -\lambda_{jik}, \lambda'_{ijk} = -\lambda''_{ikj}\). The \(SU(2)\) and color indices in Eq. (24) are contracted as follows: \(L_i Q_j D_{c_k}^c = (\nu_i d_j^c - e_i u_j^c) D_{c_k}^c\), etc. The simultaneous presence...
of all three terms in Eq. (25) will imply rapid proton decay, which can be avoided by setting the \( \lambda' = 0 \). In this case, baryon number remains an unbroken symmetry while lepton number is violated.

There are two types of to \( \beta\beta_0\nu \) decay in this model. One class dominantly mediated by heavy gluino exchange\[^{32}\] falls into the class of type II contributions discussed in the previous section. The dominant diagram of this class is shown in Fig. 3. Detailed evaluation of the nuclear matrix element for this class of models has recently been carried out by Hirsch et. al.\[^{33}\] and they have found that a very stringent bound on the following R-violating parameter can be given:

\[
\lambda'_{111} \leq 4 \times 10^{-4} \left( \frac{m_{\tilde{g}}}{100 \text{GeV}} \right)^2 \left( \frac{m_{\tilde{g}}}{100 \text{GeV}} \right)^{1/2} \tag{26}
\]

It has been recently pointed out by Faessler et al.\[^{33}\] that if one assumes the dominance of pion exchange in these processes, the limits \( \lambda'_{111} \) becomes more stringent by a factor of 2.

The second class of contributions fall into the light neutrino exchange vector-scalar type\[^{34}\] and the dominant diagram of this type is shown in Fig. 4. (where the exchanged scalar particles are the \( \tilde{s} - \tilde{s}^c \) pair). This leads to a contribution to \( \epsilon_{ee}^2 \) given by

\[
\epsilon_{ee}^2 \simeq \left( \frac{\lambda'_{131} \lambda'_{113}}{2 \sqrt{2} G_F M_b^2} \right) \left( \frac{m_b}{M_{\tilde{b}}} \right) M' \tag{27}
\]

where \( M' = (\mu \tan \beta + A_b m_0) \).

Here \( A_b, m_0 \) are supersymmetry breaking parameters, while \( \mu \) is the supersymmetric mass of the Higgs bosons. \( \tan \beta \) is the ratio of the two Higgs vacuum expectation values and lies in the range \( 1 \leq \tan \beta \leq m_t/m_b \approx 60 \). For the choice of all squark masses as well as \( \mu \) and the SUSY breaking mass parameters being of order of 100 GeV, \( A_b = 1, \tan \beta = 1 \), the following bound on R-violating couplings is obtained:

\[
\lambda'_{113} \lambda'_{131} \leq 3 \times 10^{-8} \tag{28}
\]

This bound is a more stringent limit on this parameter than the existing ones. The present limits on these parameters are \( \lambda'_{113} \leq 0.03, \lambda'_{131} \leq 0.26 \), which shows that the bound derived here from \( \beta\beta_0\nu \) is about five orders of magnitude more stringent on the product \( \lambda'_{113} \lambda'_{131} \). If the exchanged scalar particles in Fig. 9 are the \( \tilde{s} - \tilde{s}^c \) pair, one obtains a limit

\[
\lambda'_{121} \lambda'_{112} \leq 1 \times 10^{-6} \tag{29}
\]

which also is more stringent by about four orders of magnitude than the existing limits (\( \lambda'_{121} \leq 0.26, \lambda'_{112} \leq 0.03 \)).

If the quarks and leptons are composite particles, it is natural to expect excited leptons which will interact with the electron via some effective interaction involving the \( W_L \) boson. If the excited neutrino is a majorana particle, then there will be contributions to \( \beta\beta_0\nu \) decay mediated by the excited neutrinos (\( \nu^* \)). The effective interaction responsible for this is obtained from the primordial interaction:

\[
H_{\nu^*}^{\nu^*} = \frac{\lambda_{\nu^*}^{\nu^*}}{m_{\nu^*}} \bar{\nu}_{\nu^*} \gamma_{\mu \nu} (\eta_{L}^{\nu^*} \nu_L^{\nu^*} + \eta_{R}^{\nu^*} \nu_R^{\nu^*}) W_{\mu \nu} + h.c. \tag{30}
\]
Figure 4. Vector-scalar contribution in MSSM with R-parity violation.

Here L and R denote the left and right chirality states. This contribution falls into our type B heavy particle exchange category and has been studied in detail in two recent papers[35] and have led to the conclusion that it leads to a lower bound

$$m_{\nu^*} \geq 3.4 \times m_W$$

for $\lambda_W^{(\nu^*)} \geq 1$. This is a rather stringent bound on the compositeness scale.

In conclusion, neutrinoless double beta decay provides a very versatile way to probe scenarios of physics beyond the standard model. In this review, we have focussed only on the $0\nu$ mode and briefly touched on the single majoron mode. Single and multi majoron modes which test for the possibility of lepton number being a spontaneously broken global symmetry have been extensively discussed in literature[36]. The $0\nu$ mode acquires special interest in view of the recent discoveries in neutrino physics as well as certain SO(10) models predicting such spectra without contradicting the solar and atmospheric neutrino data.

REFERENCES

1. R.N. Mohapatra and P.B. Pal, “Massive Neutrinos in Physics and Astrophysics”, Second Edition, World Scientific, Singapore, 1998.
2. J. Schechter and J. W. F. Valle, Phys. Rev. D 25, 2951 (1982); E. Takasugi, Phys. Lett. 149 B, 372 (1984); B. Kayser, S. T. Petcov and P. Rosen, Private communication and see talk by B. Kayser, “Proceedings of the XXVI International Conference on High Energy Physics”, ed. J. R. Sanford, p.1153 (1992).
3. R. N. Mohapatra and G. Senjanović, Phys. Rev. Lett. 44, 912 (1980); Phys. Rev. D23, 165 (1981).
4. C. S. Aulakh and R. N. Mohapatra, Phys. Lett. 119B, 136 (1983); F. Zwirner, Phys. Lett. 132B, 103 (1983); L. Hall and M. Suzuki, Nucl. Phys. B231, 419 (1984); G. G. Ross and J. W. F. Valle, Phys. Lett. B151, 375 (1985).
5. H. Klapdor-Kleingrothaus, Prog. in Part. and Nucl. Phys., 32, 261 (1994); A. Balysh et. al., Phys. Lett. (to appear).
6. R. N. Mohapatra, in Double Beta decay and Related Topics, ed. H. Klapdor-Kleingrothaus and S. Stoica, World Scientific, 1995; p. 44; for earlier reviews see M. Doi, T. Kotani, E. Takasugi, Prog. Theor. Phys. Suppl. 83, 1 (1985); H. Primakoff and S. P. Rosen, Rep. Prog. Phys. 22, 121 (1959); W. C. Haxton and G. Stephenson, Prog. in Part. and Nucl. Phys. 12, 409 (1984); H. Grotz and H. Klapdor, The Weak Interactions in Nuclear, Particle and Astrophysics, Adam Hilger, Bristol, (1990); D. Caldwell, Nucl. Phys. Proc. Suppl. B 13, 547 (1990); M. Moe and P. Vogel, Ann. Rev. Nucl. Sc. 44, 247 (1994).
7. H. Klapdor-Kleingrothaus, these proceedings; H. Ejiri, these proceedings; F. Avignone, talk at PASCOS98 (to appear in the proceedings).
8. see the articles by P. Vogel, K. Muto, S. Stoica and S. Suhonen in Double Beta Decay and Related Topics ed. H. Klapdor-Kleingrothaus and S. Stoica, World Scientific, 1995. For a recent review, see A. Faessler and F. Simkovic, Tuebingen preprint (1998); H. Ejiri, these
proceedings.
9. A. Halprin, P. Minkowski, S. P. Rosen and H. Primakoff, Phys. Rev. D13, 2567 (1976).
10. R.N. Mohapatra and J. Vergados, Phys. Rev. Lett. 47, 1713 (1981).
11. J. Schecter and J.W.F. Valle, Phys. Rev. D25, 2951 (1982); W.C. Haxton, S.P. Rosen and G.J. Stephenson, ibid., D26, 1805 (1982); L. Wolfenstein, ibid., D26, 2507 (1982).
12. T. Kajita, these proceedings.
13. CHOOZ collaboration, M Apolonio et al, hep-ex/9711002.
14. J. Bahcall, P. Krastev and A. Smirnov, hep-ph/9807210.
15. S. P. Mikheyev and A. Smirnov, Sov. J. Nucl. Phys. 42, 913 (1985); L. Wolfenstein, Phys. Rev. D 17, 2369 (1978).
16. C. Athanassopoulos et al. Phys. Rev. Lett. 75, 2650 (1995); LSND2 C. Athanassopoulos et al. nucl-ex/9706009.
17. E. Gawiser and J. Silk, astro-ph/9806197; Science, 280, 1405 (1998).
18. D. Caldwell and R. N. Mohapatra, Phys. Rev. D 48, 3259 (1993); A. Joshipura, Z. Phys. C 64, 31 (1994).
19. F. Vissani, hep-ph/9708483.
20. R. N. Mohapatra and S. Nussinov, Phys. Lett. B 346, 75 (1995).
21. H. Minakata and O. Yasuda, Nucl. Phys. (to appear); hep-ph/9602386.
22. S. M. Bilenky, C. Giunti, C. W. Kim and M. Monteno, hep-ph/9711400.
23. H. Klapdor-Kleingrothaus, hep-ex/9802007 and these proceedings.
24. Y. Chikashige, R. N. Mohapatra and R. D. Peccei, Phys. Lett. 98B, 265 (1981).
25. M. Gell-Mann, P. Ramond and R. Slansky, in “Supergravity”, Ed. D. Freedman et al. (North-Holland, Amsterdam, 1979); T. Yanagida, Prog. Th. Phys. B135 (1978) 66; R.N. Mohapatra and G. Senjanović, Phys. Rev. Lett. 44 (1980) 912.
26. H. Georgi, S. L. Glashow and S. Nussinov, Nucl. Phys. B 193, 297 (1981).
27. R.N. Mohapatra, Phys. Rev. D34, 909 (1986).
28. M. Schwarz, Phys. Rev. D40, 1521 (1989); for a recent review, see F. Cuypers and S. Davidson, hep-ph/9609487; F. Cuypers and M. Raidal, hep-ph/9704224.
29. R. Kuchimanchi and R. N. Mohapatra, Phys. Rev. Lett. 75, 3989 (1995).
30. Z. Chacko and R. N. Mohapatra, hep-ph/9713211; C. S. Aulakh, A. Melfo and G. Senjanović, hep-ph/9707225.
31. R. N. Mohapatra, Talk at PASCOS98, (1998).
32. R. N. Mohapatra, Phys. Rev. D 34, 3457 (1986).
33. M. Hirsch, H. Klapdor-Kleingrothaus and S. Kovalenko, Phys. Rev. Lett. 75, 17 (1995); A. Faessler, S. Kovalenko, F. Simkovic and J. Schweiger, Phys. Rev. Lett. 78, 183 (1997).
34. K. S. Babu and R. N. Mohapatra, Phys. Rev. Lett. 75, 2276 (1995).
35. O. Paneva and Y. N. Srivastava, College de France Preprint, LPC 94-39; E. Takasugi, hep-ph/9506379.
36. P. Bamert, C. Burgess and R. N. Mohapatra, Nucl. Phys. B 449, 25 (1995); R. N. Mohapatra and E. Takasugi, Phys. Lett. B 211, 192 (1988); For experimental study of these processes, see J. Hellmig, M. Hirsch, H. V. Klapdor-Kleingrothaus, B. Maier and H. Pas, in Double Beta Decay and Related Topics, ed. H. V. Klapdor-Kleingrothaus and S. Stoica, World scientific (1995), p. 130.