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LEPTON PHENOMENOLOGY AND NEUTRINOLESS
DOUBLE β -DECAY IN SUPERSTRING INSPIRED THEORIES

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A B S T R A C T

The properties of neutrinos, s-leptons, s-quarks as well as the coloured isosinglet $\tilde{D}$ and $\tilde{D}^c$ are studied in superstring inspired models. Furthermore the neutrinoless double β -decay process is investigated, including light and heavy majorana neutrinos, as well as the contribution of the exotic Higgs scalars of the theory. From the current experimental limits on $0\nu\beta\beta$ -decay, useful constraints are imposed on the masses and Yukawa couplings of the theory.

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1. Introduction:

The neutrinoless double β -decay¹⁾⁻⁴⁾

$$(A, Z) \rightarrow (A, Z \pm 2) + e^{\mp} + e^{\mp} \quad (1)$$

which can occur in the presence of lepton number violating interactions is known to impose rather severe constraints on the parameters of gauge models [for a recent review see Ref. 4)]. Such are e.g. the Majorana neutrino mass, the mass of the right-handed neutrino, the couplings of Higgs scalars to leptons and even the masses of the supersymmetric partners of quarks and leptons, i.e., s-quarks and s-leptons⁵⁾. In the present paper, we will examine the neutrinoless double β -decay in the context of superstring-inspired models⁶⁾⁻¹⁴⁾. This process is ideal for such purposes since it only involves particles of the first generation.

Superstring theories⁷⁾⁻¹⁴⁾ appear to be the only candidates for a unified theory of all known interactions. The first attempt to obtain a realistic theory involved the compactification of the field theory limit of the $E_8 \times E_8$ ten-dimensional heterotic string compactified on a Calabi-Yau manifold⁷⁾. A second attempt has been made through orbifold compactifications⁹⁾. The advantage of the latter attempt is that it provides us with string theories directly with four uncompactified dimensions. Furthermore¹³⁾ there are now other ways, using the fermionic representations of directly constructing string models in four dimensions. However, although four-dimensional string theories seem to be promising, as far as we know, no model which is completely realistic¹⁴⁾ has yet been found. In fact, four-dimensional string theories do not preclude a priori the existence of phenomenological models, consistent with $SU(3) \times SU(2) \times U(1)$ theory, like those arising from the ten-dimensional heterotic string compactified on Calabi-Yau manifolds.

The field theory limit of the latter contains fundamental superfields in the adjoint 248-dimensional representation of E_8 which under $SU(3) \times E_6$ decomposes as follows:

$$248 = (1, 78) + (3, 27) + (\bar{3}, \bar{27}) + (8, 1) \quad (2)$$

The chiral fields belong to the 27-dimensional representation of E_6 which under $SO(10)$ and $SU(5)$ decomposes as follows:

$$27 = (u, d, d^c, e^+) + \nu^c + (d^c, \nu, e) + (D, \tilde{H}) + (D^c, \tilde{\tilde{H}}) + \eta \quad (3)$$

$$(16, 10) \quad (16, 1) \quad (16, \bar{5}) \quad (10, \bar{5}) \quad (10, 5) \quad (1, 1)$$

In realistic unified string theories, it is known that the compactification scale has to be of the order of the string tension¹⁰⁾. This, however, has important

constraints on superstring model building. Here we will concentrate our attention on a model which is allowed under the aforementioned restrictions. For illustrative purposes we study the model which arises when E_6 breaks down to the six-rank group:

$$SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_L \times U(1)_R \quad (4)$$

This model has L-R symmetry and is therefore related to a large class of previous GUT models⁴⁾. The decomposition of the chiral 27-dimensional representation in terms of the above subgroup chain is given in the Table.

The general superpotential which arises from the E_6 27³ invariant couplings has the form¹¹⁾⁻¹²⁾

$$W = W_0 + W_1 + W_2$$

where

$$W_0 = \lambda_1 \Phi L L^c + \lambda_2 \Phi Q Q^c + \lambda_3 \Phi \Phi \eta + \lambda_4 D D^c \eta \quad (5)$$

$$W_1 = \lambda_5 D Q Q + \lambda_6 D^c Q^c Q^c \quad (6)$$

$$W_2 = \lambda_7 D^c L Q + \lambda_8 D Q^c L^c \quad (7)$$

(generation indices have been omitted for simplicity.)

Several problems exist if all the above Yukawa couplings are present in the superpotential. If D^c is not superheavy there is the danger of too rapid proton decay via dimension five operators. This, however, can be avoided by imposing discrete symmetries¹²⁾ which forbid the existence of dangerous terms in the superpotential. In the present work, since we would like to explore the phenomenological implications of the W_2 term in the leptonic sector, we will eliminate the term W_1 by imposing a Z_2 discrete symmetry as suggested by Campbell, Ellis et al.¹²⁾

$$(3, \bar{3})_{SU(3)} \xrightarrow{Z_2} (-1) \times (3, \bar{3})_{SU(3)}$$

so that the proton does not decay too fast.

2. Lepton mixing

Before we proceed with the examination of $\nu\nu\beta\beta$ -decay we will briefly summarize the basic properties of the mass matrices following the procedure of Ref. 15). Out of the fields of the Table, those which can acquire vacuum expectation values are

$$\phi_{1/2, 1/2}, \quad \phi_{1/2, -1/2}, \quad \tilde{\nu}, \quad \tilde{\nu}^c \quad \text{and} \quad \tilde{\eta}$$

Let us begin with the leptonic sector. If one defines:

$$\begin{aligned} m_1 &= \lambda_1 \langle \phi_{1/2, 1/2} \rangle, \quad m_2 = \lambda_1 \langle \phi_{1/2, -1/2} \rangle \\ m_3 &= \lambda_1 \langle \tilde{\nu} \rangle, \quad M_1 = \lambda_1 \langle \tilde{\nu}^c \rangle, \quad M_2 = \lambda_3 \langle \tilde{\eta} \rangle \end{aligned} \quad (8)$$

one obtains the following lepton mass matrices

$$\begin{aligned} (m_e) &= \begin{array}{c|cc} & e_R^- & E_R^- \\ \hline \bar{e}_L^- & m_2 & M_1 \\ \bar{E}_L^- & m_3 & M_2 \end{array} \\ \\ (m_\nu) &= \begin{array}{c|ccccc} & \nu_R^c & N_R^c & N_R & \nu_R & \eta_R^c \\ \hline \bar{\nu}_L & 0 & 0 & M_1 & m_1 & 0 \\ \bar{N}_L & 0 & 0 & M_2 & 0 & m_1 \\ \bar{N}_L^c & M_1 & M_2 & 0 & m_3 & m_2 \\ \bar{\nu}_L^c & m_1 & 0 & m_3 & M_1' & M_{12}' \\ \bar{\eta}_L & 0 & m_1 & m_2 & M_{12}' & M_2' \end{array} \end{aligned} \quad (9)$$

The masses m_1 , m_2 and m_3 are typical Dirac masses, while M_1 and M_2 are of order 10^{11} GeV. The Majorana mass entries M_1' , M_{12}' and M_2' are required to protect the left-handed neutrino from getting a large mass, i.e., of order m_1 , a typical Dirac mass 1-10 MeV for the first generation.

They cannot arise from the superpotential of Eqs. (5)-(7) at tree level. It was, however, shown by Nandi and Sarkar¹⁵⁾ that such terms can arise by introducing dimension five operators of the type $27^2 \otimes \overline{27}^2$ of the form

$$\mathcal{L} = \frac{1}{M_c} \psi \psi \bar{S} \bar{S} \quad (10)$$

where M_c is the compactification scale ($\sim 10^{18}$ GeV) and $S, \bar{S}$ are suitable scalars. Then they showed that

$$M_1' = M_1^2/M_c, \quad M_2' = M_2^2/M_c, \quad M_{12}' = 2M_1M_2/M_c \quad (11)$$

Then the see-saw mechanism can operate and yield sufficiently light neutrinos provided that

$$|M_{12}'^2 - M_1'M_2'|^{1/2} \simeq M' \gg 10^4 \text{ GeV} \quad (12)$$

This of course means that both masses M_1 and M_2 must be superheavy ($\sim 10^{11}$ GeV). Then the low-energy group is just $SU(3)_c \times SU(2)_L \times U(1)_Y$. If, however, the low-energy group has an additional $U(1)$ factor, M_1 and M_2 cannot both be in the intermediate range and the see-saw mechanism may be doomed¹⁵⁾. Such cases have been considered in Ref. 16), but we will not consider them here.

The matrix (m_e) can be diagonalized by separate left and right unitary transformations S_L and S_R which can be obtained by diagonalizing the matrices:

$$S_L^\dagger m_e m_e^\dagger S_L = d^2, \quad S_R^\dagger m_e^\dagger m_e S_R \quad (13)$$

Thus one finds one small eigenvalue μ_1 and one superheavy μ_2 given by

$$\mu_1 \simeq (m_3 M_1 - m_2 M_2) / \sqrt{M_1^2 + M_2^2}, \quad \mu_2 \simeq \sqrt{M_1^2 + M_2^2} \quad (14)$$

The mixing matrices are

$$S_L = \begin{pmatrix} \cos \theta_L & \sin \theta_L \\ -\sin \theta_L & \cos \theta_L \end{pmatrix}, \quad S_R = \begin{pmatrix} \cos \theta_R & \sin \theta_R \\ -\sin \theta_R & \cos \theta_R \end{pmatrix} \quad (15)$$

with

$$\tan 2\theta_L = 2M_1M_2/(M_2^2 - M_1^2), \quad \tan 2\theta_R = 2 \frac{m_2M_1 + m_3M_2}{M_1^2 + M_2^2} \quad (16)$$

We immediately see that the left-handed charged leptons can be greatly admixed while the right-handed leptons remain pure.

The neutrino mass matrix can also be easily diagonalized if there are three very different mass scales in the problem, i.e., $M_1 \gg M' \gg m_1$. The mixing matrix in this case is specified in terms of two mixing angles θ_1 and θ_2 . One finds

i) One light Majorana neutrino ν_1 with mass

$$m_{\nu_1} = \xi m_1^2 / M_1'$$

$$\xi = \frac{(1 - \sin 2\theta_1 \sin 2\theta_2)(M_1' + M_2') + \cos 2\theta_1 \cos 2\theta_2 ((M_1' - M_2')^2 + 4M_{12}'^2)^{1/2}}{|M_{12}'^2 / M_1' - M_2'|} \quad (17a)$$

ii) Two intermediate mass Majorana neutrinos with masses

$$m_{\nu_2, \nu_3} = \frac{|M_1' + M_2' \pm ((M_1' - M_2')^2 + 4M_{12}'^2)^{1/2}|}{2}, (m_{\nu_2} < m_{\nu_3}) \quad (17b)$$

iii) Two degenerate superheavy neutrinos ν_4, ν_5 with masses

$$m_{\nu_4, \nu_5} = \sqrt{M_1'^2 + M_2'^2} \quad (18)$$

The current neutrino states are related to the mass eigenstates by

$$\begin{pmatrix} \nu_L \\ N_L \\ N_L^c \\ \nu_L^c \\ \eta_L \end{pmatrix} = S_L \begin{pmatrix} \nu_1 \\ \nu_2 \\ \nu_3 \\ \nu_4 \\ \nu_5 \end{pmatrix}, \quad \begin{pmatrix} \nu_R^c \\ N_R^c \\ N_R \\ \nu_R \\ \eta_R^c \end{pmatrix} = S_R \begin{pmatrix} \lambda_{CP} \nu_1 \\ \lambda_{CP} \nu_2 \\ \nu_3 \\ -\nu_4 \\ \nu_5 \end{pmatrix} \quad (19)$$

where

$$S_L = \begin{bmatrix} c_1 & 0 & 0 & \frac{s_1}{\sqrt{2}} & \frac{s_1}{\sqrt{2}} \\ -s_1 & 0 & 0 & \frac{c_1}{\sqrt{2}} & \frac{c_1}{\sqrt{2}} \\ 0 & 0 & 0 & -\frac{1}{\sqrt{2}} & -\frac{1}{\sqrt{2}} \\ \frac{m_1}{m_{\nu_2}} (c_2 c_1 - s_2 s_1) & c_2 & s_2 & 0 & 0 \\ \frac{m_1}{m_{\nu_3}} (c_2 s_1 - s_2 c_1) & -s_2 & c_2 & 0 & 0 \end{bmatrix} \quad (20)$$

where $\lambda_{CP} = \text{sign}(M'_1 M'_2 - M'^2_{12})$ is the CP-eigenvalue of the neutrinos ν_1 and ν_2 (which is expected to be negative). Also $s_i = \sin\theta_i$, $c_i = \cos\theta_i$ and $i = 1, 2$ with

$$\sin\theta_1 = \frac{M_1}{(M_1^2 + M_2^2)^{1/2}}, \quad \tan\theta_2 = \frac{2M'_{12}}{M'_2 - M'_1} = \frac{4M_1 M_2}{M_2^2 - M_1^2} = 2 \tan 2\theta_L \quad (21)$$

In the lightest neutrino case, which is a mixture of the usual weak isodoublet ν_L and isosinglet N_L , we have included the higher-order admixtures, since they are relevant in the ov $\beta\beta$ -decay. The right-handed neutrino is admixed with η to attain a mass of order $M' \sim 10^4$ GeV. The weak left-handed current can be written as:

$$j_L = 2 \bar{\nu}_L^\alpha \gamma_\lambda e_L^\alpha = 2 \cos\theta_L \left[\cos\theta_1 (\bar{\nu}_{1L} \gamma_\lambda e_L) + \frac{\sin\theta_1}{\sqrt{2}} (\nu_{4L} \gamma_\lambda e_L \bar{\nu}_{3L} \gamma_\lambda e_L) \right]_{(22)} \\ + 2 \sin\theta_L [e_L \leftrightarrow E_L]$$

where e and E now indicate the mass eigenstates. The right-handed current becomes

$$j_R = 2 \bar{\nu}_R^\alpha \gamma_\lambda e_R^\alpha = 2 \left\{ \lambda_{CP} \frac{m_1}{m_{\nu_2}} (c_2 s_2 - s_1 s_2) \bar{\nu}_{1R} \gamma_\lambda e_R + \cos\theta_2 \bar{\nu}_{2R} \gamma_\lambda e_R \right. \\ \left. + \sin\theta_2 \bar{\nu}_{3R} \gamma_\lambda e_R \right\} \quad (23)$$

From the above expression we can easily obtain the quantities which are of interest in ov $\beta\beta$ -decay⁴⁾, i.e.,

- 1) the mass of the light Majorana neutrino m_{ν_1} [see Eq. (16)]
- 2) the lepton violating parameter which involves the $j_L - j_R$ interference term

$$\eta_{RL}(k) = \frac{m_4}{m_{\nu_2}} (c_2 c_1 - s_2 s_1) k \lambda_{CP} \quad (\lambda \text{ of ref. 3}) \quad (24a)$$

$$\eta_{RL}(\varepsilon) = \frac{m_4}{m_{\nu_2}} (c_2 s_1 - s_2 c_1) \varepsilon \lambda_{CP} \quad (\eta \text{ of ref. 3}) \quad (24b)$$

where ε is the mixing of the vector bosons W_L and W_R and $k = (m_L/m_R)^2$
 3) the parameter η_R associated with right-handed interactions⁴⁾

$$\eta_R = m_p \left\langle \frac{1}{m_{\nu_R}} \right\rangle (\varepsilon^2 + k^2), \quad m_p: \text{proton mass} \quad (25)$$

and

$$\left\langle \frac{1}{m_{\nu_R}} \right\rangle = \left| \lambda_{CP} \frac{\cos^2 \theta_2}{m_{\nu_2}} + \frac{\sin^2 \theta_2}{m_{\nu_3}} \right| \quad (26)$$

Weak universality restricts the mixing angles θ_L and θ_1 so that $\cos \theta_2 \cos \theta_1 \gtrsim 0.99$ which is satisfied so long as

$$M_2 \gtrsim 10 M_1 \quad (27)$$

This in conjunction with Eq. (11) yields $\lambda_{CP} = -1$ and

$$m_{\nu_1} \approx \frac{m_4^2}{3M_1'} \left(1 - \frac{11}{2} \frac{M_1'^2}{M_2^2} \right) ; \quad \left\langle \frac{1}{m_{\nu_R}} \right\rangle \approx \frac{1}{3M_1'} \left(1 - 2 \frac{M_1'^2}{M_2^2} \right) \quad (28)$$

$$\eta_{RL} = \frac{m_4}{3M_1'} \left(1 - \frac{31}{6} \frac{M_1'^2}{M_2^2} \right) \times (k \text{ or } \varepsilon) \quad (29)$$

3. Higgs scalar contribution to $\nu\nu \beta\beta$ -decay

Additional contributions to $\beta\beta$ -decay can be induced via intermediate scalar particles. Let us first consider colourless Higgs particles. The terms in the superpotential $\lambda_2 \bar{\Phi} Q Q^c + \lambda_1 \bar{\Phi} L L^c$ can lead to the diagrams of Figs. 2a and 2b. In order to estimate such contributions, it is necessary to consider the mixing of H_2^- and H_1^+ to $\tilde{e}$ and $\tilde{e}^c$ respectively. In the absence of supersymmetry breaking their mixing and spectrum will be analogous to that of leptons estimated above. Thus one eigenstate of each charge, identified with the s-lepton, is going to be light

($m_{\tilde{e}} > 200$ GeV) while the other will be superheavy. Thus we expect very little mixing between H_1^+ and e_1^c but substantial mixing given by the angle θ_L of Eq. (16). It is thus evident that the diagram (2b) can be neglected. In the graph (2a) we encounter the isodoublet neutrino in a fashion analogous to that of Fig. 1. Thus the ratio of this contribution to the conventional one mediated by gauge bosons is

$$\chi = \left(\frac{\lambda_1 \lambda_2 \sin \theta_L}{2 G_F \mu_1^2} \right)^2 \quad (30)$$

For $\mu_1 \gtrsim 200$ GeV, $\lambda_1 = \lambda_2 = 0.1$ and $\sin \theta_L \approx 0.1$ we get $\chi < 3 \times 10^{-7}$. Additional suppression may occur due to the fact that the nuclear matrix elements encountered here are of the S-P rather than V-A type.

In the presence of $\tilde{e}^-$, the H_2^- mixing term $\lambda_2 \bar{\psi} Q Q^c$ by itself may lead to gaugino mediated $\nu \beta \beta$ -decay. Except for the mixing angle θ_L this diagram (see Figs. 3a and 3b) is similar to that of Fig. 2 of Ref. 5). Thus it gives rise to the same effective transition operator Ω_Δ , associated with S-P hadronic currents, and corresponding lepton violating parameters⁵⁾

$$\eta_{H_2^-, \tilde{\nu}} = \left(\frac{\lambda_2 \sin \theta_L \cos \theta_L}{G_F m_{\tilde{e}}^2} \right)^2 \pi \alpha \frac{m_p}{m_{\tilde{\nu}}} \quad (31a)$$

$$\eta_{H_2^-, \tilde{Z}} = \left(\frac{\lambda_2 \sin \theta_L \cos \theta_L}{G_F m_{\tilde{e}}^2} \right)^2 \frac{(-0.5 + \sin^2 \theta_W)^2}{\sin^2 2 \theta_W} \pi \alpha \frac{m_p}{m_{\tilde{Z}}} \quad (31b)$$

where $m_{\tilde{e}}$ is the mass of the lightest boson (s-electron). We shall see below that some constraints on these couplings can also arise from ordinary β -decay (see Fig. 4a).

In the presence of $\lambda_7 D^c L Q$ and $\lambda_8 D Q^c L^c$ components of the superpotential one can have new contributions to ordinary β -decay (Fig. 4b) and to $\nu \beta \beta$ -decay (Figs. 5a and 5b). For those the $\tilde{D}-\tilde{d}$ and $\tilde{D}^c-\tilde{d}^c$ mixings are necessary. The terms of the superpotential $\lambda_2 \bar{\psi} Q Q^c + \lambda_4 \bar{D} D^c \eta$ and W_2 lead to the following mass matrices

$$(m_d) = \begin{array}{c|cc} & d_R & D_R \\ \hline \bar{d}_L & m_2 & m_3 \\ \hline \bar{D}_L & M_1 & M_2 \end{array}$$

$$m_{\tilde{d}} = \begin{array}{c|ccc} & \tilde{d}^c & \tilde{D}^c & d & D \\ \hline \tilde{d}^{c*} & & & & \\ \tilde{D}^{c*} & m_d m_d^\dagger + m_{3/2}^2 & 0 & & \\ d^* & 0 & m_d^T m_d^* + m_{3/2}^2 & & \\ D^* & & & & \end{array}$$

(32)

where $m_{3/2}$ is the gravitino mass. (We have neglected the mixing due to supersymmetry breaking of d^c , D^c and d , D .)

The above matrix has the form of the transpose of the matrix (m_e) . We see that the actual values of its entry may be somewhat different than before since the relevant Yukawa couplings are different, i.e.:

$$\begin{aligned} m_2 &= \lambda_2 \langle \phi_{1/2, -1/2} \rangle, & m_3 &= \lambda_7 \langle \tilde{\nu} \rangle \\ M_1 &= \lambda_8 \langle \tilde{\nu}^c \rangle, & M_4 &= \lambda_4 \langle \tilde{\eta} \rangle \end{aligned} \quad (33)$$

The mass scales, however, are going to be similar to those of the lepton case. We thus see that we can have a sizeable $\tilde{D}^c - \tilde{d}^c$ mixing given by the mixing angle θ_2 of Eq. (16) but negligible $\tilde{D} - \tilde{d}$ mixing. In both cases we have a light particle (identical to the s-quark) with mass

$$m_{\tilde{d}}^2 = \frac{(m_3 M_1 - m_2 M_2)^2}{M_1^2 + M_2^2} + m_{3/2}^2 \quad (34)$$

and a superheavy particle with mass $M_D^2 \approx M_1^2 + M_2^2$. As we have mentioned we have a contribution to ordinary β -decay associated with Fig. 4b with amplitude

$$m_\phi = \left(\frac{\lambda_7 \sin \theta_L}{m_{\tilde{d}}} \right)^2 \bar{u}_L e_R^c \bar{\nu}_R^c d_L = k \frac{G_F}{\sqrt{2}} 2 (\bar{u}_L \gamma_5 d_L) (\bar{e}_L \gamma_5 \nu_L) \quad (35a)$$

with

$$k = \frac{1}{4\sqrt{2}} \frac{(\lambda_7 \sin \theta_L)^2}{G_F m_{\tilde{d}}^2} \quad (35b)$$

i.e., of the V-A form. From the analysis of β -decay we know that $k \lesssim 10^{-2}$, i.e.,

$$\frac{(\lambda_7 \sin \theta_L)^2}{G_F m_{\tilde{d}}^2} \lesssim 4\sqrt{2} \times 10^{-2} \quad (36)$$

The analogous process of Fig. 4a leads to S-P interaction and thus it should be more stringently constrained [see Eq. (5)].

The analysis of the contribution of the diagram (5b) is similar to that of Fig. 4 of Ref. 5). It can be parametrized in terms of an S, P, T nuclear matrix element of the operator $\mathcal{Q}(\text{SUSY})$ [see Ref. 5)] and the following set of lepton violating parameters⁵⁾.

$$\eta_{\tilde{D}^c, \tilde{g}} = \frac{4}{3} \left(\frac{\lambda_7 \sin \theta_L \cos \theta_L}{G_F m_{\tilde{d}}^2} \right)^2 \frac{4\pi\alpha_s}{16} \frac{m_D}{m_{\tilde{g}}} \quad (37a)$$

$$\eta_{\tilde{D}^c, \tilde{f}} = \frac{(\lambda_7 \sin \theta_L \cos \theta_L)^2}{(G_F m_{\tilde{d}}^2)^2} \frac{1}{9} \frac{4\pi\alpha}{16} \frac{m_P}{m_{\tilde{f}}} \quad (37b)$$

$$\eta_{\tilde{D}^c, \tilde{Z}} = \left(\frac{\lambda_7 \sin \theta_L \cos \theta_L}{G_F m_{\tilde{d}}^2} \right)^2 \frac{4\pi\alpha}{16} \left(\frac{\tan \theta_W}{6} \right)^2 \frac{m_P}{m_{\tilde{Z}}} \quad (37c)$$

where α_s is the strong coupling constant which is taken to $4\pi\alpha_s = 3.8$.

4. Constraints from $\text{ov } \beta\beta$ -decay

Neutrinoless $\beta\beta$ -decay has not yet been seen. The best limits come from the $^{76}\text{Ge} \rightarrow ^{76}\text{Se}$ and are:

$$\begin{aligned}
 T_{1/2}^{exp.}(0\nu) &> 5 \times 10^{23} y \quad (ref. 17) \\
 T_{1/2}^{exp.}(0\nu) &> 3 \times 10^{22} y \quad (ref. 18) \\
 T_{1/2}^{exp.}(0\nu) &> 5 \times 10^{22} y \quad (ref. 19)
 \end{aligned}
 \tag{38}$$

We will use the USCB-LRL data¹⁷⁾ to get the constraint

$$|\eta|^2 |ME|_{0\nu}^2 = 2.6 \times 10^{-10}$$

where $|ME|_{0\nu}^2$ are the relevant nuclear matrix elements. Using the recent nuclear matrix elements of the Ioannina-Democritos group²⁰⁾

$$\langle f | Q_\nu | i \rangle = 7.0 \quad (\text{light neutrino}) \tag{39a}$$

$$\langle f | Q_N | i \rangle = 170 \quad (\text{heavy neutrino}) \tag{39b}$$

$$\langle f | Q_A | i \rangle = 81 \quad (S-P \text{ coupling}) \tag{39c}$$

$$\langle f | Q(SUSY) | i \rangle = -246 \quad (S-P-T \text{ coupling}) \tag{39d}$$

and ignoring the interference between left- and right-handed currents we get

$$\begin{aligned}
 \left| 7.0 \frac{m_1}{m_e} + 81(\eta_{H,\tilde{p}} + \eta_{H,\tilde{g}}) - 246 \right. \\
 \left. - 246(\eta_{\tilde{D},\tilde{g}} + \eta_{\tilde{D},\tilde{t}} + \eta_{\tilde{D},\tilde{z}}) \right|^2 < 2.6 \times 10^{-10}
 \end{aligned}
 \tag{40a}$$

$$|170 \eta_R|^2 < 2.6 \times 10^{-10} \tag{40b}$$

Under the simplifying assumption of one mechanism at a time being operative we get

$$| \langle m_{\nu_1} \rangle | \leq 1.1 \text{ eV} \quad \text{or} \quad \xi \frac{m_h^2}{M_1} < 1.1 \text{ eV} \tag{41a}$$

$$\eta_R \leq 9.2 \times 10^{-8} \quad \text{or} \quad (\epsilon^2 + k^2) m_p \left\langle \frac{1}{m_{\nu_R}} \right\rangle \leq 9.2 \times 10^{-8} \tag{41b}$$

The best limits on the parameter η_{RL} come from Ref. 21), i.e.,

$$\eta_{LR}(\epsilon) \leq 3.6 \times 10^{-6}, \quad \eta_{LR}(k) \leq 2.8 \times 10^{-8} \quad (42)$$

In the special case of the approximation of Eq. (28) (i.e., $\xi = 1/3$) Eqs. (41a) and (41b) become

$$\frac{m_1^2}{3M_1'} \leq 1.1 \text{ eV}, \quad \frac{m_P}{3M_1'} (k^2 + \epsilon^2) \leq 2.9 \times 10^{-8} \quad (43)$$

Assuming now that $k \approx 0.05$ and $\epsilon \approx 0.01$ we obtain

$$M_1' \geq 0.9 \times 10^4 \text{ GeV} \quad \text{and} \quad m_1 \leq 5.5 \text{ MeV} \quad (44)$$

These in turn imply

$$\eta_{RL}(k) \leq 1.0 \times 10^{-8}, \quad \eta_{RL}(\epsilon) \leq 2.0 \times 10^{-9} \quad (45)$$

which are within the bounds of Eq. (42). From Eq. (40) we get

$$|\eta_{\tilde{D}^c, \tilde{g}} + \eta_{\tilde{D}^c, \tilde{p}} + \eta_{\tilde{D}^c, \tilde{z}}| \leq 6.5 \times 10^{-8} \quad (46)$$

This limit is quite stringent and forces us to assume that

$$\frac{(\lambda_f \sin \theta_L)^2}{G_F m_{\tilde{g}}^2} = 4\sqrt{2} \times 10^{-5} \quad (47)$$

i.e., three orders of magnitude below the limit of Eq. (36) (otherwise the s-quarks would have to be much heavier). Then the Eq. (41) becomes

$$\left(1.8 \frac{m_P}{m_{\tilde{g}}} + 3.6 \times 10^{-3} \frac{m_P}{m_{\tilde{p}}} + 2.7 \times 10^{-4} \frac{m_P}{m_{\tilde{z}}}\right) \left(\frac{m_P}{m_{\tilde{d}}}\right)^2 \leq 6.5 \times 10^{-8} \quad (48)$$

Clearly the last term is negligible and the second term is unimportant since we do not expect $m_{\tilde{g}} > 500 m_{\tilde{Y}}$. Thus assuming that $m_{\tilde{g}} \approx 100 m_P$ we get

$$m_{\tilde{d}} \gtrsim 500 \text{ GeV} \quad (49)$$

Substituting this in Eq. (47) and noting that $\sin\theta_L \lesssim 0.1$ we get

$$\lambda_7 \approx 1.3 \times 10^{-3} \quad (50)$$

Returning to the diagram (3a) and assuming that

$$\frac{\lambda_1 \lambda_2 \sin\theta_L}{G_F m_{\tilde{d}}^2} = 4\sqrt{2} \quad (51)$$

in conjunction with the limit of Eq. (40) we obtain

$$|\eta_{H;\tilde{\nu}} + \eta_{H;\tilde{e}}| < 2 \times 10^{-7} \quad (52)$$

which, processing as above, can be translated into limits

$$m_{\tilde{e}} \gtrsim 360 \text{ GeV}, \quad \lambda_1 \approx \lambda_2 \approx 1.3 \times 10^{-3} \quad (53)$$

if we take $m_{\tilde{\gamma}} \approx 4m_p$

Furthermore, if we assume that the above estimates are generation-blind, we get from the limits (41b), (46) and (47) that the branching ratio^{4),22)} for (μ^-, e^+) is

$$R(\mu^-, e^+) \lesssim 10^{-29} - 10^{-28} \quad (54)$$

i.e., substantially larger than the prediction of traditional SO(10) model but still essentially unobservable.

5. Conclusions

In conclusion we have shown that the present experimental limits on $0\nu\beta\beta$ -decay can provide useful constraints on the parameters of superstring-inspired models. In particular the following estimates have been made.

Using successful constraints of the standard model, like weak universality, we have been able at first to obtain restrictions on the mixing angles of the left-handed current in Eq. (22). We have found further that the contribution of the ordinary Higgs scalars is suppressed compared to that of gauge bosons [see Eq. (30)]. However, a substantial mixing of D^C - d^C is found to occur, which leads to stringent constraints on the Yukawa coupling λ_7 which is of order 10^{-3} [Eq. (50)], through the D^C - d^C -quark mixing mediated graph of Fig. 5a. The same result holds also for the couplings λ_1, λ_2 as one can check in Eq. (53). Such constraints are apparently very stringent since here we are forced to accept much smaller Yukawa couplings than the usual ones of the standard model.

Nevertheless, these small values should be compared with those which arise from the flavour non-conserving processes and other rare reactions which are investigated in similar superstring-inspired models in Ref. 12).

Finally the following bounds have emerged on the mass spectrum:

- i) The Dirac mass m_l related to the neutrino has been found to be less than 5.5 MeV [Eq. (44)];
- ii) The mass of the squark $m_{\tilde{q}^C}$ is found to be at least 500 GeV [Eq. (49)];
- iii) The mass of the slepton $m_{\tilde{e}}$ is at least 360 GeV [Eq. (53)].

In the above estimates we have assumed that $m_{\tilde{g}} \gtrsim 500$ GeV, $m_{\tilde{\gamma}} \approx 5$ GeV. Admittedly there are uncertainties in such models. Further, certain simplified assumptions were necessary. They are, however, reasonable and in line with assumptions commonly made in gauge theories.

One may therefore view our predictions, especially Eqs. (45), (49) and (53), with a reasonable degree of confidence.

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Fermions	Name (symbol) Bosons	SO(10) dim.	SU _c (3) dim.	SU _L (2) I _L	SU _R (2) I _R	Y _L	Y _R
$Q = \begin{pmatrix} u \\ d \end{pmatrix}_L$	$\tilde{Q} = \begin{pmatrix} \tilde{u} \\ \tilde{d} \end{pmatrix}$	16	3	1/2	0	1/3	0
$Q^c = \begin{pmatrix} d^c \\ -u^c \end{pmatrix}_L$	$\tilde{Q}^c = \begin{pmatrix} \tilde{d}^c \\ -\tilde{u}^c \end{pmatrix}$		$\bar{3}$	0	1/2	0	-1/3
$L = \begin{pmatrix} \nu \\ e \end{pmatrix}_L$	$\tilde{L} = \begin{pmatrix} \tilde{\nu} \\ \tilde{e} \end{pmatrix}$		1	1/2	0	-1/3	-2/3
$L^c = \begin{pmatrix} e^c \\ -\nu^c \end{pmatrix}_L$	$\tilde{L}^c = \begin{pmatrix} \tilde{e}^c \\ -\tilde{\nu}^c \end{pmatrix}$		1	0	1/2	2/3	1/3
D_L	$\tilde{D}$	10	3	0	0	-2/3	0
D_L^c	$\tilde{D}^c$		$\bar{3}$	0	0	0	2/3
$\tilde{\Phi} = \begin{pmatrix} \tilde{H}_1 & \tilde{H}_1^+ \\ \tilde{H}_2^- & \tilde{H}_2^+ \end{pmatrix}$	$\Phi = \begin{pmatrix} H_1^+ & H_1^+ \\ H_2^- & H_2^+ \end{pmatrix}$		1	1/2	1/2	-1/3	1/3
η_L	$\tilde{\eta}$	1	1	0	0	2/3	-2/3

TABLE

The decomposition of the chiral 27-dimensional E_6 representation according to $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_L \times U(1)_R$. The SO(10) quantum numbers are also indicated^{*}.

^{*}) Sometimes we will find it convenient to use the notation $(\tilde{H}_1^+, \tilde{H}_1^+) = (N^c, E^c)$ and $(\tilde{H}_2^+, \tilde{H}_2^-) = (N, E)$ used in Ref. 15) with the corresponding I_{3L} , I_{3R} assignments.

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FIGURE CAPTIONS

Figure 1: Diagram leading to conventional $0\nu\beta\beta$ -decay.

Figure 2: Possible mechanisms leading to $0\nu\beta\beta$ -decay mediated by Higgs particles provided that they have light components.

Figure 3: A process analogous to that of Fig. 2 involving gauginos rather than neutrinos provided that there is substantial $H_2^-, \tilde{e}^-$ and e^c, H_1^+ mixing. The latter is negligible in our model.

Figure 4: New contributions to ordinary β -decay induced by $H_2^-(a)$ and $\tilde{D}^c(b)$ provided that they have substantial amount of light components.

Figure 5: New contributions to $0\nu\beta\beta$ -decay arising from D^c, d^c mixing (a) and D, d mixing (b). In our model the latter is negligible.

Fig. 1

Fig. 2a

Fig. 2b

Fig. 3a

Fig. 4a

Fig. 3b

Fig. 4b

Fig. 5a

Fig. 5b