Observable Physics from Superstring Exotic Particles: Small Dirac Neutrino Masses

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Certain superstring theories contain new exotic particles and couplings active at low energies. We show that they offer a solution to a long-standing problem of simultaneously having naturally small Dirac neutrino masses and vanishing Majorana masses. It is therefore consistent with (a) lack of observation of neutrinoless double β decay, (b) possible evidence for nonvanishing neutrino mass, and (c) a recently proposed solution to the solar neutrino problem.

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Superstring theories[1] represent the most promising candidate for a consistent theory of all fundamental interactions including gravity. In spite of major unsolved questions already at the string level, some general implications for the resulting four-dimensional physics have already been worked out.[2-4]

In the most appealing superstring theory based on the $E_8 \otimes E_8$ heterotic string, matter fields belong to the **27** representation of E_6:

$$\mathbf{27} = \{ Q \equiv \binom{u}{d},\ u^c,\ d^c,\ D^c,\ D,\ L \equiv \binom{\nu}{e},\ e^c,\ \nu^c,\ N,\ H \equiv \binom{H^+}{H^0},\ \bar{H} \equiv \binom{\bar{H}^0}{H^-} \}, \qquad (1)$$

where $(D^c D^c D^c H^- \bar{H}^0) + (DDDH^0 H^+)$ form the $5 \oplus 5^*$ representation of SU(5) and ν^c and N are SU(5) superfield singlets. In addition to the couplings of the standard model, the low-energy theory obtained from the superstring may contain the couplings[5]

$$\lambda_L HL\nu^c + \lambda H\bar{H}N + hDD^cN + f_L + f_Q, \qquad (2)$$

where $f_L = \lambda_1 D^c LQ + \lambda_2 De^c u^c + \lambda_3 Dd^c \nu^c$ and $f_Q = \lambda_4 DQQ + \lambda_5 D^c u^c d^c$. Notice that the extra terms in (2) are the only ones allowed by the underlying E_6 symmetry. Moreover, we want to emphasize that in spite of the presence of the E_6 gauge grand unification at some scale close to the compatification scale, the Yukawa couplings of the matter fields do not respect the E_6 Clebsch-Gordan relations. This is a precious clue for avoiding phenomenological disasters like a rapid proton decay or large neutrino masses. Moreover we avoid typical SU(5)-type quark-lepton mass relations.

In a general class of dimensional reduction schemes, the four-dimensional theory presents a residual supersymmetry and a no-scale structure.[5,6] In this framework all energy scales are determined dynamically[5] and particles contained in (1) acquire masses below a teraelectronvolt. Consequently the simultaneous presence of all the couplings in (2) leads to unacceptable phenomenological consequences. Indeed, the term $\lambda_1 HL\nu^c$ gives large neutrino masses as well as large mixings $\propto \lambda_L \langle \nu^c \rangle$ between leptons and Higgs fermions with a consequent large nonconservation of lepton number. The presence of both f_L and f_Q terms leads to unbearably fast baryon-nonconserving phenomena. However, f_L and f_Q conserve B separately, with the choices $B(D) = +\frac{1}{3}$, $B(D^c) = -\frac{1}{3}$ and $B(D) = -\frac{2}{3}$, $B(D^c) = +\frac{2}{3}$. Thus we must forbid the λ_L term and f_L or f_Q in (2). The most appealing way is to exclude some terms in the superstring formalism through topological considerations. Unfortunately, we do not yet know how to achieve this with Calabi-Yau manifolds; it is, anyway, encouraging that the choice $\lambda_L = 0$ and $f_L = 0$ or $f_Q = 0$ can be implemented by imposition of discrete symmetries.

The imposition of this discrete symmetry leading to the B and L conservation in the low-energy sector, apart from avoiding too fast a proton decay, prevents any dilution of the hard-won cosmological baryon asymmetry. In fact, B nonconservation is not forbidden at higher mass scales since we have enforced the conservation of B only in the light sector. Finally, notice that by imposing this discrete symmetry we are creating a connection between apparently different problems, such as B nonconservation, flavor-changing neutral currents, and the weak universality which would have been destroyed by the D-d mixing.[4]

A general analysis of flavor-changing neutral currents, $g-2$, d_e^n, CP nonconservation, etc. in these theories is being considered elsewhere.[7] Here we focus our attention on the case $\lambda_L = 0$, $f_L \neq 0$, and $f_Q = 0$ and discuss the limits imposed on the parameters by the experimental bounds of some rare processes, such as $K^+ \to \pi^+ \nu\nu$, $\pi^0 \to \mu e$, μ capture, and $\mu \to e\gamma$. Using the range of coupling constants thus determined, we propose a solution to the long-standing problem of simultaneously having naturally small Dirac masses for neutrinos and vanishing Majorana contributions.

The tree-level $\tilde{D}$- and $\tilde{D}^c$-mediated diagrams give rise to the reactions

$$d_L\bar{d}_L \rightarrow \nu_L\bar{\nu}_L \quad (d_L\bar{s}_L \rightarrow \nu_L\bar{\nu}_L), \quad d_R\bar{u}_R \rightarrow \bar{\nu}_R e_R,$$
$$u_L\bar{u}_L \rightarrow e_L\bar{e}_L, \quad d_L\bar{d}_R \rightarrow \nu_R\bar{\nu}_L,$$
$$d_R\bar{d}_R \rightarrow \nu_R\bar{\nu}_R, \quad u_L\bar{u}_R \rightarrow e_R\bar{e}_L, \tag{3}$$
$$u_R\bar{u}_R \rightarrow e_R\bar{e}_R, \quad d_L\bar{u}_R \rightarrow e_R\bar{\nu}_L, \quad u_L\bar{d}_R \rightarrow \nu_R\bar{e}_L.$$

For simplicity we restricted ourselves to the reactions involving first-generation fermions. The presence of analogous reactions which mix generations is understood. The reactions in (3) lead to a bunch of interesting rare procesess: $\pi_0 \rightarrow \nu_L\bar{\nu}_R$, $K^+ \rightarrow \pi^+\nu\bar{\nu}$, $K_L \rightarrow \pi^0\nu\bar{\nu}$, $K_L \rightarrow \nu\bar{\nu}$, $\pi^0 \rightarrow e^+e^-$, $\pi^0 \rightarrow \mu e$, $D^0 \rightarrow \mu e + X$, $\psi \rightarrow \mu e + X$, μ + nucleus $\rightarrow e$ + nucleus, at the tree level. At the one-loop level the most interesting process is $\mu \rightarrow e\gamma$.

At the tree level, the most stringent bounds on λ_1 and λ_3 come from the processes $K^+ \rightarrow \pi^+\nu\nu$ and μ + nucleus $\rightarrow e$ + nucleus. Taking for definiteness the entries of the λ_1 and λ_3 matrices to be of the same order, for $K^+ \rightarrow \pi^+\nu\nu$ (Fig. 1) we obtain

$$\frac{\Gamma(K^+ \rightarrow \pi^+\nu\bar{\nu})}{\Gamma(K^+ \rightarrow \pi^0 e^+\nu)} \simeq \frac{\lambda_1^4}{g^4}\left(\frac{m_W}{m_{\tilde{D}}}\right)^4, \tag{4}$$

where $m_{\tilde{D}}$ denotes the mass of $\tilde{D}$ or $\tilde{D}^c$. From the experimental bound,[11]

$$\Gamma(K^+ \rightarrow \pi^+\nu\bar{\nu})/\Gamma(K^+ \rightarrow \pi^0 e^+\nu) < 0.33\times10^{-5},$$

we infer

$$\lambda_1 < \frac{m_{\tilde{D}}}{m_W} 8\times(10^{-3}\text{–}10^{-2}), \tag{5}$$

for $m_{\tilde{D}}$ in the 100–200-GeV range. The diagrams which are relevant for the μ-e conversion are depicted in Fig. 2. Summing up three contributions, we finally get[8]

$$\frac{\Gamma(\mu+\text{nucleus} \rightarrow e+\text{nucleus})}{\Gamma(\mu+\text{nucleus} \rightarrow \nu+\text{nucleus})} \simeq \frac{\frac{1}{8}[(Z+A)/Z](1/e^2 m_{\tilde{D}}^2)^2}{16G_F^2}[\lambda_1^2+\lambda_2^2+\lambda_2\lambda_1]^2, \tag{6}$$

where we take, throughout this paper, the DD^c mixing term to be $O(m_{\tilde{D}}^2)$.

Taking the experimental bound[9] for Eq. (6) coming from $\mu^- + \mathrm{S}_{32} \rightarrow e^- + \mathrm{S}_{32}$ to be 10^{-11} and $\lambda_1 \sim \lambda_2 = \lambda$ we obtain

$$\frac{\lambda_1}{m_{\tilde{D}}} < \frac{5\times10^{-6}}{1\ \text{GeV}} \rightarrow \lambda_1 < 10^{-3}. \tag{7}$$

Note that the process $\pi^0 \rightarrow \mu e$ which is obtainable again through tree-level diagrams does not yield any appreciable bound on the λ couplings since

$$\frac{\Gamma(\pi^0 \rightarrow \mu e)}{\Gamma(\pi^0 \rightarrow X)} \sim \frac{\lambda_1^4}{g^4}\left(\frac{m_W}{m_{\tilde{D}}}\right)^4 10^{-8}, \tag{8}$$

and the experimental bound[9] on $\Gamma(\pi^0 \rightarrow \mu e)/\Gamma(\pi^0 \rightarrow X)$ is 7×10^{-8}. The structure of our couplings does not allow for a tree-level $K \rightarrow \mu e$ or $K^+ \rightarrow \pi^+\mu e$ decay.

At the one-loop level, the most relevant process to provide a bound on the λ couplings is certainly $\mu \rightarrow e\gamma$ (Fig. 3). Evaluating the graph in Fig. 3, we obtain the branching ratio

$$B(\mu \rightarrow e\gamma) \lesssim \frac{3\alpha}{\pi}\frac{1}{G_F^2 m_{\tilde{D}}^4}\left(\frac{m_b}{m_\mu}\right)^2. \tag{9}$$

From (19), and the experimental value[9] $B(\mu \rightarrow e\gamma) \leq 1.7\times10^{-10}$, we get

$$\frac{\lambda_1}{m_{\tilde{D}}} < \frac{10^{-5}}{1\ \text{GeV}} \rightarrow \lambda_1 < 10^{-3}, \tag{10}$$

in the same range as the bound derived from μ capture.

Thus, we find that for $m_{\tilde{D}} \simeq m_{\tilde{D}^c} \simeq$ 100–200 GeV the couplings λ_1, λ_2, and λ_3 can be of the typical order of the Yukawa couplings of the second fermionic generation ($< 10^{-3}$) without any conflict with the present experimental bounds.

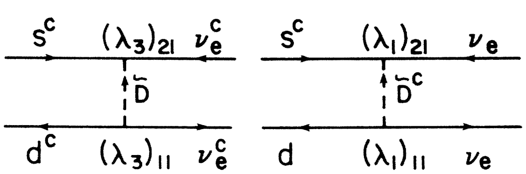


FIG. 1. Diagrams contributing to $K^+ \rightarrow \pi^+\nu\nu$.

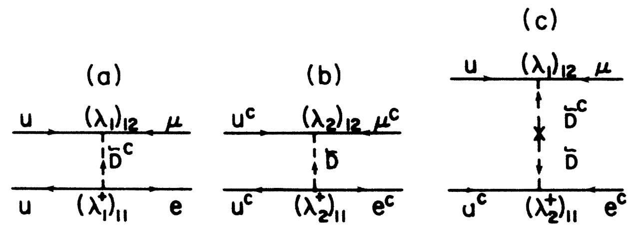


FIG. 2. Diagrams contributing to μ + nucleus $\rightarrow e$ + nucleus.

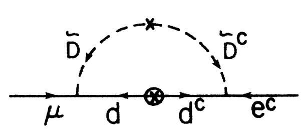


FIG. 3. Diagram contributing to $\mu \rightarrow e\gamma$.

The presence of the f_L couplings allows for the radiative origin of nonvanishing Dirac masses for neutrinos (for characteristic diagram see Fig. 4):

$$m_\nu \sim (1/4\pi^2)\lambda_1\lambda_3 M_d, \tag{11}$$

where M_d denotes the down-quark mass matrix. If we use the upper bound $\lambda \sim 10^{-3}$ and take, for instance, diagonal λ matrices, (11) yields

$$m_{\nu_\tau} \sim 50 \text{ eV}, \quad m_{\nu_\mu} \sim 1 \text{ eV}, \\ m_{\nu_e} \sim 10^{-1} \text{ eV}. \tag{12}$$

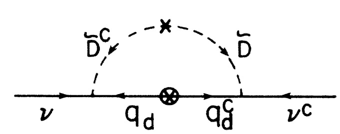
Hence, without any particular fine tuning on λ's, we succeed in obtaining small neutrino Dirac masses which are compatible with the cosmological bounds. The enforcement of lepton-number conservation forbids any Majorana entry in the neutrino mass matrix.

The possibility of obtaining naturally small Dirac masses for neutrinos is quite interesting from the phenomenological point of view and, indeed, several efforts in this direction were made in the past.[10] The Dirac nature of neutrino masses automatically forbids neutrinoless double-beta decay in agreement with the present experimental situation. If the Lyubimov experiment[11] with evidence of a ν_e mass in the 20–40-eV range is confirmed, this can be accommodated in our scheme provided that a large intergenerational mixing in λ matrices is present. Here we would like to focus on another issue which has received much attention quite recently. It has been pointed out that the mixing of different neutrino species while they propagate through matter leads to a possible solution of the solar neutrino problem.[12] For this to occur, the mass difference $m_{\nu_\mu}^2 - m_{\nu_e}^2 \equiv \Delta$ must be $\sim 6\times 10^{-5}$ eV^2. In our model a rough estimate of Δ yields

$$\Delta \simeq [(1/4\pi^2)\lambda_1\lambda_3]^2 m_s^2, \tag{13}$$

where m_s is the mass of the s quark. Here we took the intergenerational mixing in λ to be small. Hence if we want our mechanism to generate neutrino masses to be compatible with the Mikheyev-Smirnov[12] proposal we might set $\lambda_1 \sim \lambda_3 \sim 0.4\times 10^{-4}$. Remembering that the Yukawa coupling of the quark s, for instance, is $\sim 10^{-4}$, we conclude that $\Delta \sim 10^{-5}$ eV^2 can be obtained in our model without a particular fine tuning.

Going back to the rare processes that we have previously analyzed, taking $\lambda_1 \sim 10^{-4}$ yields branching ratios some 4 orders of magnitude below the present experimental upper bounds for $\mu \rightarrow e\gamma$ and μ conversion. This should not discourage the search for such processes since the rates go as λ^4 and are quite sensitive to our ignorance of the structure of the γ matrices. Finally, the presence of extra particles in the low-energy regime allows for a mechanism of radiative production of fermionic masses. The smallness of these contributions tells us that they can be relevant in discussion of the properties of the fermions of the lightest generation.



q_d = d,s,b

FIG. 4. Characteristic diagram giving rise to neutrino Dirac mass contributions.

In conclusion, we have proposed here a mechanism to generate small neutrino Dirac masses in the context of low-energy supersymmetric models inspired by the $E_8 \otimes E_8$ heterotic string theory. These masses are compatible with the cosmological bounds and, indeed, it is possible to implement the Mikheyev-Smirnov proposal for the solar neutrino problem without any particular fine tuning of the parameters.

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