RADIATIVE NEUTRINO DECAY IN SUPERSTRING MODELSK. Enqvist^{*)}Research Institute for Theoretical Physics
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ABSTRACT

We calculate the rate for the radiative decay of neutrinos in E_6 superstring models with light colour triplet isosinglets h and $\bar{h}$. In contrast to the standard model, the rate is not suppressed either by the GIM mechanism or by helicity matching. Comparison with astrophysical data allows us to set an upper limit on the Yukawa couplings $\bar{\lambda}QL\bar{h}$ and $\lambda d^c \nu^c h$. Taking $M_{h,\bar{h}} = O(100 \text{ GeV})$ we find $\lambda_{e,\mu}, \bar{\lambda}_{e,\mu} < 10^{-3}-10^{-4}$, compatible with m_{ν_e}, m_{ν_μ} in the range 10-100 eV. From the cosmological requirement of not photo-ionizing light elements after nucleosynthesis, it follows that the ν_τ must lie in the same mass range and consequently $\lambda_\tau, \bar{\lambda}_\tau < 10^{-3}$.

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The heterotic superstring theory^{1),2)} yields, after (2,2) compactification of the six surplus dimensions to a Calabi-Yau manifold with $SU(3)$ holonomy, an E_6 gauge theory coupled to $N = 1$ supergravity^{3),4)}. Breaking of E_6 by the Wilson-loop mechanism^{3),5)} necessarily implies a low-energy gauge symmetry larger than the standard $SU(3)_c \times SU(2)_L \times U(1)_Y$. For example, in the no-scale model⁶⁾, which arises as the effective $d = 4$ supergravity theory of superstrings after simple truncation of the ten-dimensional theory⁷⁾, the low-energy symmetry is

$$G = SU(3)_c \times SU(2)_L \times U(1) \times U(1)' \quad (1)$$

In this model⁶⁾ there are no intermediate mass scales between the Planck scale $M_P \approx 10^{19}$ GeV and the dynamically-induced scale $\mu \approx 10^2 - 10^3$ GeV, where first the extra $U(1)'$ invariance and then the standard electroweak symmetry $SU(2)_L \times U(1)_Y$ are broken. This structure of the symmetry breaking implies that the whole matter multiplets, having quantum numbers of a fundamental 27 representation of E_6 , survive with all their important phenomenological consequences to low energies.

One phenomenological aspect of the superstring theories, quite essential for model building, is the question of neutrino mass. In E_6 superstring models fermions can acquire mass, at the tree level, only from the vacuum expectation values of 27 multiplets. As a consequence, the standard seesaw mechanism⁸⁾ where the right-handed neutrino obtains a very large Majorana mass is not possible in these models. Several scenarios have recently been suggested to resolve this problem. They include, e.g., models where a small mass is induced to the observed neutrinos by an intermediate mass scale through higher-order terms in the Lagrangian⁹⁾, or through a conspiring mixing pattern among neutral fermions¹⁰⁾ (there are altogether five neutral states in a 27; see the classification below), or where the job is done by E_6 -singlet zero mass fields¹¹⁾ which appear in (2,0) compactifications¹²⁾. These scenarios, however, do not seem phenomenologically viable¹³⁾. It is also possible that some mirror particles, i.e., the particles living in 27^* , remain massless down to the scale μ ; in that case, one can try¹⁴⁾ to solve the neutrino mass problem due to mixing of neutrinos and mirror neutrinos with no need of an intermediate scale. However, there is the difficulty that in superstring models there are no renormalizable couplings between 27 and 27^* .

In the no-scale superstring models⁶⁾, none of the above mechanisms are conceivable. These models do not develop intermediate scales¹⁵⁾, and absence of

lepton-number-violating processes seems to require massless neutrinos at the tree level. This situation can follow in superstring theories as a result of certain discrete symmetries originating, e.g., in the topological properties of the internal manifold. Although the well-known nucleosynthesis limit on light neutrinos¹⁶⁾ may need some revision¹⁷⁾, there remains the uncomfortable possibility of having too many light neutrinos.

If, however, the colour triplet isosinglet fields, leptoquarks, belonging to the representation 10 in the $SO(10)$ decomposition $27 = 16 + 10 + 1$ remain light, as they will in the no-scale models, a Dirac mass which is tiny but still lying in an interesting range can be naturally generated through loop corrections¹⁸⁾, as shown in Fig. 1.

In the present paper we will study in detail a question closely related to the neutrino mass problem, namely radiative stability of neutrinos in the E_6 superstring model. We will consider loop diagrams of a similar type to those generating neutrino masses which will allow, as is pointed out also in Ref. 20), for a radiative decay $\nu \rightarrow \nu' \gamma$ of a neutrino many orders of magnitude more rapid than that in the standard model (see Fig. 2). The reason for the rapidity of this decay in the superstring model is the fact that diagrams involving leptoquarks, in contrast to the ones mediated by the standard electroweak gauge interactions, are not helicity-suppressed, which yields an enhancement of m_d/m_ν in the amplitude. We will confront the ensuing radiative lifetime $\tau(\nu \rightarrow \nu' \gamma)$ with the existing astrophysical and cosmological constraints to derive bounds for the leptoquark Yukawa couplings.

Let us start by exhibiting the most general cubic 27^3 superpotential which describes the interactions among the various components of the supermultiplet 27:

$$\begin{aligned} P = & L H e^c + Q \bar{H} d^c + Q H u^c + L H \nu^c \\ & + H \bar{H} N + h \bar{h} N + Q Q h + u^c d^c \bar{h} \\ & + Q L \bar{h} + d^c \nu^c h + u^c e^c h . \end{aligned}$$

(2)

The fields appearing here have their transformation properties under the group $G = SU(3) \times SU(2) \times U(1)_Y \times U(1)'$ as follows:

$$\begin{aligned}
 Q &= \begin{pmatrix} u \\ d \end{pmatrix} = (3, 2)_{\frac{1}{6}, \frac{1}{3}} \quad , \quad u^c = (3^*, 1)_{-\frac{2}{3}, \frac{1}{3}} \quad , \quad d^c = (3^*, 1)_{\frac{1}{3}, -\frac{1}{6}} \quad , \\
 L &= \begin{pmatrix} \nu \\ e^- \end{pmatrix} = (1, 2)_{-\frac{1}{2}, -\frac{1}{6}} \quad , \quad e^c = (1, 1)_{1, \frac{1}{3}} \quad , \quad \nu^c = (1, 1)_{0, \frac{5}{6}} \quad , \\
 h &= (3, 1)_{-\frac{1}{3}, \frac{2}{3}} \quad , \quad \bar{h} = (3^*, 1)_{\frac{1}{3}, -\frac{1}{6}} \quad , \quad H = \begin{pmatrix} H^+ \\ H^- \end{pmatrix} = (1, 2)_{-\frac{1}{2}, \frac{1}{6}} \quad , \\
 \bar{H} &= \begin{pmatrix} H^+ \\ H^- \end{pmatrix} = (1, 2)_{\frac{1}{2}, -\frac{2}{3}} \quad , \quad N = (1, 1)_{0, \frac{5}{2}} \quad . \quad (3)
 \end{aligned}$$

All the couplings given in Eq. (2) cannot obviously be non-vanishing; a simultaneous appearance of the last five terms would contradict the observed stability of the proton since the colour triplet fields h and $\bar{h}$ are assumed to be light. It is a specific virtue of the superstring models (not shared by the conventional GUTs) that one may assume the diquark couplings of h and $\bar{h}$, $QQh + u^c d^c \bar{h}$, to vanish while allowing their leptoquark couplings $QL\bar{h} + d^c \nu^c h + u^c e^c h$ and the rest of the superpotential (3) to survive¹⁵⁾.

The part of the superpotential (2) relevant for the neutrino mass and radiative decay is, in a more complete notation, given by

$$\bar{\lambda}_{ij}^k \varepsilon_{\alpha\beta} Q_i^\alpha L_j^\beta \bar{h}^k + \lambda_{ij}^k d_i^c \nu_j^c h^k + \text{h.c.} \quad , \quad (4)$$

where $i, j, k = 1, \dots, n_G$ are flavour indices and $\alpha, \beta = 1, 2$ are SU(2) indices. The Yukawa couplings $\bar{\lambda}$ and λ are determined by the properties of the internal manifold and, in the absence of detailed knowledge of that manifold, must be taken as free parameters. There are, of course, many phenomenological constraints on these couplings, in particular on $\bar{\lambda}$'s, since the light leptoquark exchanges will contribute to different low-energy processes. These have been studied in Refs. 18) to 20). Assuming $M_{h, \bar{h}} \approx 100$ GeV, one obtains generically $\bar{\lambda} \lesssim 10^{-2} - 10^{-3}$. In general, the diagonal couplings are less severely constrained than the non-diagonal ones.

For the decay $\nu \rightarrow \nu' \gamma$, non-vanishing non-diagonal Yukawa couplings are not necessary, as long as there is flavour mixing in the down quark and leptoquark $h, \bar{h}$ systems. In the following, we will for simplicity assume that the couplings

really are diagonal and denote $(\bar{\lambda})_{ij}^k \equiv (\bar{\lambda})_i^k$.

Let us call U , V and $\bar{V}$ the matrices which give the down quarks, leptoquarks h and leptoquarks $\bar{h}$, respectively, in terms of the corresponding mass eigenstates. We can then derive the following expressions for the amplitudes of the two diagrams given in Figs. 2a and 2b:

$$\begin{aligned}
 A_a(\nu_j \rightarrow \nu_i \gamma) &= -i\left(\frac{e}{3}\right) \epsilon_{\mu} \sum_{r,s} \bar{\lambda}_j \lambda_i \bar{V}_{jn} V_{in} U_{sj} U_{is}^{\dagger} \int \frac{d^4 k}{(2\pi)^4} \left[\bar{u}_j(p-q)(1-\gamma_5) \frac{i}{k-q-m_n} \right. \\
 &\quad \left. \times \gamma^{\mu} \frac{i}{k-m_n} (1-\gamma_5) u_i(p) \frac{i}{(k-p)^2-M_s^2} \right] \\
 &= -\frac{i}{16\pi^2} \left(\frac{e}{3}\right) \sum_{n,s} \frac{m_n}{M_s^2} F_a\left(\frac{m_n}{M_s}\right) \bar{\lambda}_j \lambda_i \bar{V}_{jn} V_{in} U_{sj} U_{is}^{\dagger} \bar{u}_j(p-q) \sigma^{\mu\nu} \epsilon_{\mu} q_{\nu} (1-\gamma_5) u_i(p), \quad (5) \\
 A_b(\nu_j \rightarrow \nu_i \gamma) &= -i\left(\frac{e}{3}\right) \epsilon_{\mu} \sum_{n,s} \bar{\lambda}_j \lambda_i \bar{V}_{jn} V_{in} U_{sj} U_{is}^{\dagger} \int \frac{d^4 k}{(2\pi)^4} \left[\bar{u}_j(p-q)(1-\gamma_5) \frac{i}{k-q-m_r} \right. \\
 &\quad \left. \times (1-\gamma_5) u_i(p) \frac{i}{k^2-M_s^2} \frac{i(2k+q)^{\mu}}{(k+q)^2-M_s^2} \right] \\
 &= -\frac{i}{16\pi^2} \left(\frac{e}{3}\right) \sum_{n,s} \frac{m_n}{M_s^2} F_b\left(\frac{m_n}{M_s}\right) \bar{\lambda}_j \lambda_i \bar{V}_{jn} V_{in} U_{sj} U_{is}^{\dagger} \bar{u}_j(p-q) \sigma^{\mu\nu} \epsilon_{\mu} q_{\nu} (1-\gamma_5) u_i(p),
 \end{aligned}$$

where ϵ and q are the polarization vector and the momentum of the photon and p is the momentum of the decaying neutrino ν_j . The functions F_a and F_b are for small values of the argument $\xi = m_n/M_s$ given by

$$\begin{aligned}
 F_a(\xi) &= \frac{3}{2} + \frac{5}{2}\xi + (1+3\xi)\ln\xi, \\
 F_b(\xi) &= 1 + 3\xi + 2\xi\ln\xi.
 \end{aligned} \quad (6)$$

A striking fact is that the amplitudes are proportional to the down-quark masses, rather than to the neutrino masses, which would be the case in the standard model²¹⁾. This yields a huge enhancement factor^{*)} and the question arises whether the result can still be made compatible, within any sensible range of parameters, with the astrophysical constraints on the radiative neutrino decay.

There are many unknown parameters on which the rate $\Gamma(\nu_j \rightarrow \nu_i \gamma)$ depends, i.e., the masses M of the leptoquarks, the leptoquark Yukawa couplings $\lambda, \bar{\lambda}$ and the flavour-mixing parameters of the leptoquark system. Parameters on which we do have experimental information are the down-quark masses and mixings, the latter

*) A similar enhancement factor may arise in the standard model extended to include right-handed interactions; see Ref. 22).

being given in the Kobayashi-Maskawa matrix [the matrix U in Eq. (5)]. The lepto-quark masses originate from soft supersymmetry breaking $h\bar{h}N$ couplings¹⁵⁾; they cannot thus exceed the scale $M \approx 10^2 - 10^3$ GeV, otherwise the original motivation for supersymmetry would be lost. For the flavour mixing in the $h, \bar{h}$ sector, we may, in the absence of any better knowledge, let the phenomenologically-observed mixing among the ordinary down quarks guide us, and assume that this is of the same order of magnitude.

So that we may obtain an order-of-magnitude estimate, we now consider the decay $\nu_\mu \rightarrow \nu_e \gamma$ with the simplifying assumptions that only two down quarks, d and s , and two leptoquarks, h_1 and h_2 , are relevant. We further assume that the h_1, h_2 mixing is parametrized by an angle θ_h (which we will take equal to the familiar Cabibbo angle θ_c). Taking only the leading terms in the ratios $m_{d,s}/M_{1,2}$ into account, the amplitude is given by

$$\mathcal{A}_{\text{arb}}(\nu_\mu \rightarrow \nu_e \gamma) = \frac{i}{128\pi^2} \left(\frac{e}{3}\right) \bar{\lambda}_\mu \lambda_e \sin 2\theta_c \sin 2\theta_h \frac{(m_d - m_s)(M_1^2 - M_2^2)}{M_1^2 M_2^2} \\ \times \bar{u}_{\nu_\mu}(p-q) \sigma^{\mu\nu} \epsilon_\nu q_\nu (1-\gamma_5) u_{\nu_e}(p). \quad (7)$$

Note that the GIM cancellation does not occur with full strength here, since the amplitude is proportional to the masses of the loop particles.

The amplitude (7) corresponds to the lifetime τ , given by

$$\tau(\nu_\mu \rightarrow \nu_e \gamma) = \frac{\alpha (\bar{\lambda}_\mu \lambda_e)^2}{9\pi^4 2^4} \sin^2 2\theta_c \sin^2 2\theta_h \frac{(m_d - m_s)^2 (M_1^2 - M_2^2)^2}{(M_1 M_2)^4} \frac{1}{E_0^3}, \quad (8)$$

where

$$E_0 = \frac{m_{\nu_\mu}^2 - m_{\nu_e}^2}{2m_{\nu_\mu}} \quad (9)$$

is the photon energy. For a numerical estimate we take $\sin \theta_h = \sin \theta_c = 0.23$, $m_s - m_d \approx 0.2$ GeV, $M_1 + M_2 = 10^2$ GeV and $|M_1 - M_2| \approx 10$ GeV, with which the lifetime becomes

$$\tau^{-1} \approx (0.6 \times 10^{12} \text{ sec})^{-1} (\lambda_e \bar{\lambda})^2 \left(\frac{E_0}{25 \text{ eV}} \right)^3. \quad (10)$$

Let us recall that the radiative lifetime in the standard model is $\tau > 10^{35}$ sec (for $E_0 = 25$ eV)²¹⁾.

There are several observational lower limits for the radiative lifetime of neutrinos from astrophysics²³⁾. Assuming the neutrino abundance derived from the standard Big Bang cosmology, the most stringent limit for the energy range $E_0 = 5-13$ eV (i.e., $m_{\nu_\mu} \approx 10-27$ eV, if $m_{\nu_e} = 0$), $\tau \gtrsim 10^{23}$ sec, is obtained from the measurement of the far ultra-violet spectrum of the background radiation²⁴⁾. Above the Lyman limit $E_0 = 13.6$ eV an even more severe limit, $\tau \gtrsim 10^{24}$ sec, is implied by the fact that there exist neutral hydrogen clouds in galactic haloes²⁵⁾; for shorter decay times these would have been ionized by the photon flux^{*)}.

Let us confront these observational limits with our theoretical result (10). Allowing E_0 to get values in the range $E_0 = 5-50$ eV, for example, leads to the following upper limit for the Yukawa couplings:

$$\lambda_e \bar{\lambda} \lesssim 7 \times 10^{-6} - 7 \times 10^{-8}, \quad (11)$$

or generically

$$\lambda, \bar{\lambda} \lesssim 10^{-3} - 10^{-4}. \quad (12)$$

It must be stressed that the bounds (11) and (12) depend quite considerably on the various assumptions adopted. For example, allowing masses an order of magnitude higher for the leptoquarks ($M_1 + M_2 \approx 1$ TeV, $|M_1 - M_2| \approx 0.1$ TeV, say) would permit the Yukawa coupling constants $\lambda, \bar{\lambda}$ to be of the size of a typical gauge coupling constant. On the other hand, assuming $E_0 > 50$ eV would have an effect working into the opposite direction. It is also possible that the level of mass

*) Observation of no enhanced radiation from the directions of certain galaxy clusters pushes the lower limit of τ still higher in some narrow bands of the photon energy E_0 . For example, observations on the Virgo cluster indicate²⁶⁾ $\tau \gtrsim 0.6 \times 10^{25}$ sec for $E_0 = 8-10$ eV.

degeneracy among the leptoquarks is different than assumed here ($|M_1 - M_2|/(M_1 + M_2) \approx 10\%$), or that the size of the flavour mixing in the leptoquark sector is considerably larger or smaller than that among the ordinary down quarks given by the KM matrix.

The above analysis is valid for a general class of E_6 models with light leptoquarks. In the special case where the loop diagrams depicted in Fig. 1 are the only source of neutrino masses, there will be a dynamical relationship between these masses and the radiative decay amplitude, a fact which will affect our conclusions. From Fig. 1 we obtain generically^{18),20)}

$$m_\nu \approx \frac{\lambda^2}{16\pi^2} m_f \quad (13)$$

(with $\lambda \sim 10^{-3}$ this gives $m_{\nu_e} \approx 1$ eV, $m_{\nu_\mu} \approx 4$ eV and $m_{\nu_\tau} \approx 60$ eV) and if we assume that all neutrino masses are of the same order of magnitude, we can write

$$\Gamma(\nu \rightarrow \nu' \gamma) \approx \frac{\alpha m_\nu^5}{M_{h,h}^4} \sin^4 \theta \times O(10^{-2}) \quad (14)$$

where $O(10^{-2})$ takes into account numerical factors and $\sin^4 \theta$ the flavour mixing. [Incidentally, Eqs. (13) and (14) are true also in the case where h-quarks propagate in the loop, because the trilinear scalar coupling $\tilde{d}\tilde{d}^c H$ is suppressed by the d-quark mass.] The width (14) has the same general form as the corresponding width in the standard model; only the W mass is replaced by $M_{h,h}$. Thus if we assume neutrinos in the 10-100 eV mass range (which means $\lambda \sim 10^{-3}$), the radiative lifetime will, as we know from the standard case²³⁾, pass the astrophysical limits by several orders of magnitude.

One could still consider the possibility that the tau neutrino ν_τ may become much heavier than the ν_e and ν_μ due to the fact that the Yukawa couplings $\lambda_\tau, \bar{\lambda}_\tau$ are not as strongly constrained by the laboratory experiments as the couplings for the other two families. Equation (13) tells us that masses up to $m_{\nu_\tau} \approx 60$ MeV are possible if $\lambda_\tau, \bar{\lambda}_\tau \lesssim O(1)$. The radiative lifetime for such a neutrino, $\tau \approx 10^4$ sec, would be short enough to avoid any contradiction with the constraints from the primordial nucleosynthesis and the energy density of the Universe, as well as from the astrophysical arguments considered above. However, the decay photons would dissociate the primordial deuterium and helium²⁷⁾ which have the

photofission thresholds at 2.2 MeV and 20 MeV respectively. To escape this contradiction one should require $\tau \lesssim 50 \text{ sec}^{27)*}$, which is quite impossible to achieve in the present case. We can thus conclude that ν_τ should lie in the electron volt region and that consequently the bound $\lambda_\tau, \bar{\lambda}_\tau \lesssim 10^{-3}$ should hold.

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*) It has recently been argued²⁸⁾ that this limit may actually be considerably higher, viz. $\tau \lesssim 5 \times 10^3 \text{ sec}$. If so, there would be a marginal possibility to have $m_{\nu_\tau} \approx 70 \text{ MeV}$, which, however, may soon become ruled out by laboratory measurement of the mass.

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FIGURE CAPTIONS

- Fig. 1 : Characteristic diagram generating a Dirac mass for the neutrino in the E_6 superstring model. The h and $\bar{h}$ are colour triplet isosinglet scalars with their mass in the range 10^2 - 10^3 GeV. The supersymmetric counterpart would be the diagram where $\tilde{d}^c$ squarks and $\tilde{h}, \bar{\tilde{h}}$ quarks propagate in the loop.
- Fig. 2 : Diagrams for radiative neutrino decay $\nu_j \rightarrow \nu_i \gamma$. Comparing decay rate to astrophysical data allows us to derive upper limits for the Yukawa couplings λ and $\bar{\lambda}$. The indices i, j, r, s indicate flavour.

Fig. 1

Fig. 2