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Neutrino Masses in Superstring Derived Standard–like Models

ALON E. FARAGGI*

SSC Laboratory, 2550 Beckleymeade Ave., Dallas TX 75237

and

Department of Physics, Weizmann Institute of Science
Rehovot 76100, Israel‡

and

Edi Halyo†

Department of Physics, Weizmann Institute of Science
Rehovot 76100, Israel

ABSTRACT

We propose a new scenario in a class of superstring derived standard–like models that explains the suppression of the left–handed neutrino masses. Due to nonrenormalizable terms and the breaking of the $U(1)_{Z'}$ symmetry a generalized see–saw mechanism takes place. Contrary to the traditional see–saw mechanism in GUTs, the see–saw scale and the right–handed neutrino mass scale are suppressed relative to the $U(1)_{Z'}$ breaking scale.

* e–mail address: fhalon@weizmann.bitnet
‡ Permanent address
† e–mail address: jphalyo@weizmann.bitnet

1. Introduction

Superstring theories [1] are believed to provide a viable framework for the consistent unification of gravity with the gauge interactions. The superstring unification scale is at the Planck scale while at the electro–weak scale the Standard Model is in remarkable agreement with experimental observations. Thus, an extremely important task is to try to derive the Standard Model from the superstring theories.

In ref. [2,3,4], realistic superstring standard–like models were constructed in the free fermionic formulation [5] with the following characteristics:

1. Three and only three generations of chiral fermions. There are no additional generations and mirror generations that presumably get massive at a high scale.
2. The gauge group is $SU(3) \times SU(2) \times U(1)_{B-L} \times U(1)_Y \times U(1)^n \times hidden$.
3. There are enough scalar doublets and singlets to break the symmetry in a realistic way and to generate a realistic fermion mass hierarchy.
4. Proton decay from dimension four operators is supressed due to the gauged $B-L$ symmetry. There is a stringy doublet-triplet splitting mechanism that projects out the color triplets and leaves the Higgs doublets in the spectrum. Thus, proton decay from dimension five operators is expected to be supressed [4].
5. String gauge coupling unification can be obtained in these models due to the presence of additional color triplets in vector–like representations from exotic sectors [6].
6. These models suggest an explanation for the top–bottom quark mass hierarchy and for the generation mass hierarchy in general. At the trilinear level of the superpotential only the top quark gets a nonvanishing mass. The mass terms for the bottom quark and for the lighter quarks and leptons are obtained from nonrenormalizable terms. Thus, only the top mass is characterized by the electroweak scale and the masses for the lighter quarks and leptons are naturally supressed [3,4].

In this paper we examine the problem of neutrino masses in the superstring standard–like models. Like many other extensions of the Standard Model the superstring standard–like models contain right-handed neutrinos. At trilevel an underlying $SO(10)$ symmetry dictates that for every up–like quark mass term there is a neutrino Dirac mass term, with $\lambda_\nu = \lambda_u$. Thus, the τ neutrino mass is equal to the top quark mass. Clearly a disasterous relation, unless there is a see–saw type mechanism. In this paper we demonstrate that a see–saw mechanism is present in the superstring standard–like models.

We study the renormalizable and nonrenormalizable contributions to the neutrino mass matrix. We show that a large see–saw in the neutrino mass matrix is generated by condensates of a non–Abelian hidden gauge group. As a result the left–handed neutrino masses are supressed whereas the right–handed neutrino masses are pushed above the TeV scale. However, we find a new feature in our superstring derived see–saw mechanism unprecedented in traditional Grand Unified Theories (GUTs). In traditional superstring inspired GUTs the see–saw scale is proportional to the GUT scale. For example in the flipped $SU(5)$ the see–saw matrix takes the form [7],

$$(\nu_i \quad \nu_j^C \quad \phi_m)\begin{pmatrix} 0 & (km_U)_{ij} & 0 \\ (km_U)_{ij} & 0 & M_G \\ 0 & M_G & O(M_W) \end{pmatrix}\begin{pmatrix} \nu_i \\ \nu_j^C \\ \phi_m \end{pmatrix}, \tag{1}$$

where m_u is the up quark mass matrix, M_G is the scale of $SU(5)\times U(1)$ breaking. The mass eigenvalues are $m_\nu \sim \langle\phi\rangle(\frac{km_u}{M_G})^2$ and $m_{N,\phi} \sim M_G$. Thus the mass scale of the right–handed neurtrinos is the GUT scale. Similarly in the superstring inspired standard–like model [8], the see–saw scale and the mass of the right–handed neutrinos are proportional to the scale of $U(1)_{Z'}$ breaking.

In the superstring derived see–saw mechanism, the $U(1)_{Z'}$ is broken by hidden sector condensates. We show that although the scale of the hidden sector condensates is high, $\Lambda_h \sim 10^{14}\ GeV$, the masses of the right–handed neutrinos are suppressed relative to the $U(1)_{Z'}$ breaking scale.

2. The superstring model

To illustrate the superstring induced see–saw mechanism, we examine the model of Ref. [2]. This model is constructed in the free fermionic formulation [5], with a set of eight boundary condition vectors for all the world–sheet fermions. The first five vectors in the basis consist of the NAHE* set, $\{1, S, b_1, b_2, b_3\}$ [9,4]. In addition to the first five vectors, the basis contains three additional vectors. These vectors and the choice of generalized GSO projection coefficients are given in Table 1, where the notation of Ref. [4] is used.

The gauge group after application of the generalized GSO projections is

Observable† : $SU(3)_C \times U(1)_C \times SU(2)_L \times U(1)_L \times U(1)^6$

Hidden# : $SU(5)_H \times SU(3)_H \times U(1)^2$.

The weak hypercharge is uniquely given by $U(1)_Y = \frac{1}{3}U(1)_C + \frac{1}{2}U(1)_L$. The orthogonal combination is given by $U(1)_{Z'} = U(1)_C - U(1)_L$. In the observable sector there are six horizontal $U(1)$ symmetries. The first three, $U(1)_{r_j}$ ($j = 1,2,3$), correspond to the right–moving world–sheet currents $\bar{\eta}_1\bar{\eta}_1^*$, $\bar{\eta}_2\bar{\eta}_2^*$ and $\bar{\eta}_3\bar{\eta}_3^*$. The last three, $U(1)_{r_{j+3}}$ ($j = 1,2,3$), correspond to the right–moving world–sheet currents, $\bar{y}^3\bar{y}^6$, $\bar{y}^1\bar{\omega}^5$ and $\bar{\omega}^2\bar{\omega}^4$ respectively. For every right–moving $U(1)$ symmetry there is a left–moving global $U(1)$ symmetry. The first three correspond to the charges of the supersymmetry generator χ^{12}, χ^{34} and χ^{56}. The last three, $U(1)_{\ell_{j+3}}$ ($j = 1,2,3$), correspond to the complexified left–moving fermions y^3y^6, $y^1\omega^5$ and $\omega^2\omega^4$. Finally the model contains six Ising model sigma operators which are obtained by pairing a left–moving real fermion with a right–moving real fermion, $\sigma^i_\pm = \{\omega^1\bar{\omega}^1, y^2\bar{y}^2, \omega^3\bar{\omega}^3, y^4\bar{y}^4, y^5\bar{y}^5, \omega^6\bar{\omega}^6\}_\pm$.

The full massless spectrum and cubic level superpotential were presented in Ref. [2]. Here we list only the states which are relevant for the neutrino see–

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* This set was first constructed by Nanopoulos, Antoniadis, Hagelin and Ellis (NAHE) in the construction of the flipped $SU(5)$ [9]. *nahe*=pretty, in Hebrew.

† $U(1)_C = \frac{1}{2}U(1)_{B-L}$, $U(1)_L = \frac{1}{2}U(1)_{T_{3_R}}$.

\# Hidden here means that the states which are identified with the chiral generations do not transform under the hidden gauge group.

saw mass matrix. The sectors b_1, b_2 and b_3 produce three generations of chiral fermions, G_1, G_2 and G_3, with horizontal symmetries. For every generation, G_j there are two right–moving, $U(1)_{r_j}$ and $U(1)_{r_{j+3}}$, symmetries. For every right–moving $U(1)$ gauged symmetry, there is a corresponding left–moving global $U(1)$ symmetry, $U(1)_{\ell_j}$ and $U(1)_{\ell_{j+3}}$. Each sector b_1, b_2 and b_3 has two Ising model operators, $(\sigma_4,\ \sigma_5)$, $(\sigma_2,\ \sigma_6)$ and $(\sigma_1,\ \sigma_3)$, respectively, obtained by pairing a left–handed real fermion with a right–handed real fermion. The pairings of real fermions, which produce the Ising model operators, is noted in Table 1. The nonvanishing cubic level and higher order terms in the superpotential must be invariant under all the symmetries of the superstring model. The Ising model operators must produce nonvanishing corralators after all picture changing operations have been performed.

The Neveu–Schwarz sector produces three pairs of electroweak doublets h_1, $\bar{h}_1$, h_2, $\bar{h}_2$ and h_3, $\bar{h}_3$. Each pair carries $(U(1)_{\ell_j}; U(1)_{r_j})$ charges, $(j = 1,2,3)$, respectively. The Neveu–Schwarz sector also produces three pairs of $SO(10)$ singlets Φ_{12}, $\bar{\Phi}_{12}$, Φ_{13}, $\bar{\Phi}_{13}$, Φ_{23}, $\bar{\Phi}_{23}$ with $U(1)_{r_j}$ charges. Finally, the Neveu–Schwartz sector gives rise to three singlet states that are neutral under all the U(1) symmetries. $\xi_{1,2,3}:\ \ \chi^{12}_{\frac{1}{2}}\bar{\omega}^3_{\frac{1}{2}}\bar{\omega}^6_{\frac{1}{2}}|0\rangle_0,\ \chi^{34}_{\frac{1}{2}}\bar{y}^5_{\frac{1}{2}}\bar{\omega}^1_{\frac{1}{2}}|0\rangle_0,\ \chi^{56}_{\frac{1}{2}}\bar{y}^2_{\frac{1}{2}}\bar{y}^4_{\frac{1}{2}}|0\rangle_0$.

The $S + b_1 + b_2 + \alpha + \beta$ sector gives

$$h_{45} \equiv [(1,0);(2,1)]_{-\frac{1}{2},-\frac{1}{2},0,0,0,0} \qquad D_{45} \equiv [(3,-1);(1,0)]_{-\frac{1}{2},-\frac{1}{2},0,0,0,0} \quad (2a,b)$$
$$\Phi_{45} \equiv [(1,0);(1,0)]_{-\frac{1}{2},-\frac{1}{2},-1,0,0,0} \qquad \Phi_1^{\pm} \equiv [(1,0);(1,0)]_{-\frac{1}{2},\frac{1}{2},0,\pm1,0,0} \quad (2c,d)$$
$$\Phi_2^{\pm} \equiv [(1,0);(1,0)]_{-\frac{1}{2},\frac{1}{2},0,0,\pm1,0} \qquad \Phi_3^{\pm} \equiv [(1,0);(1,0)]_{-\frac{1}{2},\frac{1}{2},0,0,0,\pm1} \quad (2e,f)$$

(and their conjugates $\bar{h}_{45}$, etc.). The states are obtained by acting on the vacuum with the fermionic oscillators $\bar{\psi}^{4,5}, \bar{\psi}^{1,\ldots,3}, \bar{\eta}^3, \bar{y}^3 \pm i\bar{y}^6, \bar{y}^1 \pm i\bar{\omega}^5, \bar{\omega}^2 \pm i\bar{\omega}^4$, respectively (and their complex conjugates for $\bar{h}_{45}$, etc.).

The sectors $b_i + 2\gamma + (I)$ $(i = 1,..,3)$ give vector representations which are $SU(3)_C \times SU(2)_L \times U(1)_L \times U(1)_C$ singlets (see Table 2). These states transform

in the fundamental representation under the hidden non–Abelian gauge groups. The sectors $b_1 + b_3 + \alpha + \gamma + (I)$ and $b_2 + b_3 + \beta + \gamma + (I)$ produce states with vanishing weak hypercharge and with non vanishing $U(1)_{Z'}$ charge (see Table 3). In particular the sector $b_1 + b_3 + \alpha + \gamma + (I)$ produces 5 and $\bar{5}$ of the hidden $SU(5)$ gauge group with $U(1)_{Z'}$ charge and vanishing weak hypercharge. However, the entire massless spectrum does not contain 5 and $\bar{5}$ of the hidden $SU(5)$ group that carry $U(1)_Y$ charge.

The standard–like models contain an anomalous $U_A(1)$ symmetry which induces a D–term and destabilizes the vacuum [10]. To preserve supersymmetry at the Planck scale, one must satisfy the F and D constraints arising from the superpotential by giving VEVs to the scalar fields [11]. We will consider the solution to the constraints given by Eq. 12 of Ref. [2]:

$$|\Phi_{45}|^2 = 3\Delta_{13} = 3|\bar{\Phi}_1^-|^2 = 3|\bar{\Phi}_2^-|^2 = 3|\Phi_3^+|^2 = \frac{g^2}{16\pi^2}\frac{1}{\sqrt{2\alpha'}} \tag{3}$$

where $\Delta_{13} = |\bar{\Phi}_{13}|^2 - |\Phi_{13}|^2$ and $\sqrt{2\alpha'} = 2\sqrt{8\pi}/gM_{Pl}$. We would like to emphasize that the solution, Eq. (3), preserves F flatness to all orders of nonrenormalizable terms. For solutions that break $U(1)_{Z'}$ close to the Planck scale, higher order terms are expected to violate F flatness [13]. We therefore assume that VEVs which break $U(1)_{Z'}$ are suppressed relative to the $SO(10)$ singlet VEVs.

3. Neutrino mass matrix

The cubic level and higher order terms in the superpotential are derived by applying the rules of Ref. [12]. Order N nonrenormalizable terms are of the form $cffh(\Phi/M)^{N-3}$ where f, h and Φ are fermions, scalar doublets and singlets respectively. M is a Planck scale mass to be defined later. The coefficients c can be calculated from string amplitudes and are $\sim O(1)$. For our choice of boundary condition vectors given in Table 1, at the cubic level, there are Dirac mass terms

only for the $+\frac{2}{3}$ charged quarks and for the neutrinos, with $\lambda_u = \lambda_\nu$,

$$(u^c_{L_1} Q_1 \bar{h}_1 + N^c_{L_1} L_1 \bar{h}_1 + u^c_{L_2} Q_2 \bar{h}_2 + N^c_{L_2} L_2 \bar{h}_2 + u^c_{L_3} Q_3 \bar{h}_3 + N^c_{L_3} L_3 \bar{h}_3). \quad (4)$$

At the cubic level of the superpotential only two pairs of Higgs doublets which are combinations of h_1,h_2, h_{45} and $\bar{h}_1$,$\bar{h}_2$,$\bar{h}_{45}$ remain light. At an intermediate energy scale one additional pair becomes heavy [13]. For a specific choice of VEVs, the light doublets will be $\bar{h}_2$ and h_{45} [13]. Therefore, there is a cubic level mass term only for the top quark and for the τ neutrino (the heaviest generation is the one with index 2), with $\lambda_t = \lambda_{\nu_\tau}$. The μ neutrino obtains a Dirac mass term at order $N = 5$.

$$N_2 L_2 (\bar{h}_1 \bar{\Phi}^+_i \bar{\Phi}^-_i + \bar{h}_{45} \bar{\Phi}_{23} \Phi_{45}) \quad (5a)$$

$$N_1 L_1 (\bar{h}_2 \Phi^+_i \Phi^-_i + \bar{h}_{45} \bar{\Phi}_{13} \Phi_{45}) \quad (5b)$$

Since the VEVs of the scalar fields are generally $\sim O(10^{-1})M$ the Dirac mass term for the μ neutrino is suppresed relative to the τ neutrino by $\left(\frac{\langle\phi\rangle}{M}\right)^2 \sim 0.01$. As shown in Ref. [14] a Dirac mass term for the electron neutrino can only arise from higher order terms which contain VEVs that break $U(1)_{Z'}$. VEVs that break $U(1)_{Z'}$ are expected to be suppressed relative to the $SO(10)$ singlet VEVs [13].

Mixing between the generations is induced by the states from the sectors $b_j+2\gamma$ [15]. However, the mixing terms are suppressed relative to the leading diagonal terms and therefore their effect is small. We neglect them in this paper.

In the superstring derived standard–like model the states in the see–saw mass matrix are the neutral states in the left–handed doublets L_1, L_2, L_3, the right–handed neutrinos N_1, N_2, N_3, from the sectors b_1, b_2 and b_3 and the $SO(10)$ singlets $\Phi = \{\Phi_{13}, \bar{\Phi}_{13}, \Phi^+_1, \bar{\Phi}^-_1, \Phi^+_2, \Phi^+_3, \bar{\Phi}^-_3\}$ from the Neveu–Schwarz sector and from the sector $S + b_1 + b_2 + \alpha + \beta$. These $SO(10)$ singlets remain massless after

the symmetry breakdown for the choice of VEVs given by Eq. (3) and therefore appear in the low energy spectrum. There are other light singlets but only these mix with the right–handed neutrinos. A nonvanishing VEV for the right–handed neutrino will induce rapid proton decay from dimension four operators [13]. Therefore, the right–handed neutrinos are assumed to have vanishing VEVs.

The neutrino–singlet mixing terms arise from higher order nonrenormalizable terms, which are obtained according to the rules of Ref. [12]. The general form of the terms is fixed by the symmetries of the model. The coefficients of these terms can in principle be calculated from string scattering amplitudes, and are of order one [12]. The nonrenormalizable terms contributing to the see–saw mass matrix appear at orders $N = 6$, $N = 7$ and $N = 8$.

At order $N = 6$ there is only one see–saw term for N_2:

$$N_2 \Phi_2^+ H_{19} \bar{T}_2 H_{23} \Phi_{45} \tag{6}$$

At order $N = 7$ we obtain the following see–saw terms for N_3:

$$N_3 \bar{\Phi}_3^- H_{19} \bar{T}_3 H_{25} \Phi_{45} \bar{\Phi}_{13} \tag{7a}$$

$$N_3 \Phi_3^+ H_{19} \bar{T}_3 H_{23} \Phi_{45} \xi_2 \tag{7b}$$

Other see–saw terms that appear at this order are suppressed relative to the ones above since they contain more VEVs that break $U(1)_{Z'}$. For N_1 the mixing terms are at order $N = 8$:

$$N_1 \Phi_1^+ H_{19} \bar{T}_1 H_{23} \Phi_{45} \xi_2 \xi_2 \tag{8a}$$

$$N_1 \Phi_{45} H_{19} \bar{T}_1 H_{25} \Phi_1^- \bar{\Phi}_{13} \xi_2 \tag{8b}$$

$$N_1 \bar{\Phi}_{13} H_{19} \bar{T}_1 H_{25} \bar{\Phi}_1^- \Phi_{45} \xi_2 \tag{8c}$$

$$N_1 \bar{\Phi}_{13} H_{19} \bar{T}_1 H_{25} \bar{\Phi}_1^- \Phi_{45} \xi_2 \tag{8d}$$

$$N_1 \bar{\Phi}_1^- H_{24} \bar{T}_1 H_{25} \Phi_{45} \bar{\Phi}_{13} \xi_2 \tag{8e}$$

The see–saw terms that are suppressed relative to those above are neglected. Up to order $N = 8$ there are no mass terms for the $SO(10)$ singlets for our specific choice of VEVs in Eq. (3). At order $N = 9$, there is a mass term of the form $\Phi_2^+\Phi_2^+\bar{T}_2\bar{T}_2T_2T_2\Phi_{45}\Phi_{45}\xi_2M^{-6}$ for Φ_2^+. This term contains four VEVs that break $U(1)_{Z'}$. Other singlets get masses from higher order ($N > 9$) terms.

In Eqs. (6,7,8) the $5, \bar{5}$ states, H_{19} and $\bar{T}$ of the hidden $SU(5)$ group form condensates. These condensates carry nonzero $U(1)_{Z'}$ charges and therefore will break $U(1)_{Z'}$. In addition we make the assumption that additional singlets (H_{23}, H_{25}, H_{26}) with nonzero $U(1)_{Z'}$ charge get a VEV at the same scale. The $SO(10)$ singlets from the Neveu–Shwarz sector and the sector $S + b_1 + b_2 + \alpha + \beta$ obtain VEVs by the cancellation of the anomalous D–term equation. As a result the terms in Eqs. (5,6,7) induce the effective see–saw terms of the form

$$N\Phi\left(\frac{\Lambda_{Z'}}{M}\right)^3\left(\frac{\langle\phi\rangle}{M}\right)^n M \tag{9}$$

where $\Lambda_{Z'}$ is the scale at which $SU(5)$ condensates form and break $U(1)_{Z'}$ and $M = M_P/2\sqrt{8\pi} \sim 10^{18}$ GeV. n is 1,2 and 3 for N_2, N_3 and N_1 respectively. We take the $SO(10)$ singlet mass terms, m_ϕ to be generically of the form

$$\Phi\Phi\left(\frac{\Lambda_{Z'}}{M}\right)^4\left(\frac{\langle\phi\rangle}{M}\right)^m M \tag{10}$$

where $m = 3$ for Φ_2^+ and $m > 3$ for the other singlets. It should be emphasized that for the see–saw mechanism to operate there must exist singlets that couple to the right–handed neutrinos as in Eq. (9) and that remain massless at trilevel for the flat directions. If such singlets exist then the higher order mass terms for these singlets will be generically of the form given by Eq. (10).

The neutrino mass matrix therefore takes the following form for each generation in the basis (ν_L, N^C, Φ)

$$\begin{pmatrix} 0 & km_u & 0 \\ km_u & 0 & m_\chi \\ 0 & m_\chi & m_\phi \end{pmatrix} \tag{11}$$

with $m_\chi \sim \left(\frac{\Lambda_{Z'}}{M}\right)^3 \left(\frac{\langle\phi\rangle}{M}\right)^n M$ and $m_\phi \sim \left(\frac{\Lambda_{Z'}}{M}\right)^4 \left(\frac{\langle\phi\rangle}{M}\right)^m M$. The mass eigenstates are mainly ν, N and ϕ with a small mixing and with the eigenvalues

$$m_\nu \sim m_\phi \left(\frac{k m_u}{m_\chi}\right)^2 \qquad m_N, M_\phi \sim m_\chi \tag{12}$$

The constant k gives the effects of Yukawa coupling renormalization. At M $\lambda_u = \lambda_\nu$. These two couplings run differently since up–quarks and neutrinos have different quantum numbers. In order to find their ratio at the weak scale,M_W, one has to solve the renormalization group equations for these two couplings which are

$$\frac{d\lambda_1^t}{dt} = \frac{\lambda_1^t}{32\pi^2}\left(12\lambda_1^{t2} + 2\lambda_2^{b2} + 2\lambda_4^2 + 2\lambda_7^2 - 4\left(\frac{8}{3}g_3^2 + \frac{3}{2}g_2^2 + \frac{13}{30}g_1^2\right)\right) \tag{13}$$

$$\frac{d\lambda_4}{dt} = \frac{\lambda_4}{8\pi^2}\left(2\lambda_4^2 + \frac{3}{2}\lambda_1^{t2} + \frac{1}{2}\lambda_3^{\tau 2} + \frac{1}{2}\lambda_7^2 - \left(\frac{3}{2}g_2^2 + \frac{3}{10}g_1^2\right)\right) \tag{14}$$

where $\lambda_1^t, \lambda_2^b, \lambda_3^\tau, \lambda_4, \lambda_7$ are the Yukawa couplings of the top, bottom,τ, τ neutrino and scalar Higgs respectively. g_1, g_2 and g_3 are the $U(1)_Y$, $SU(2)_L$ and $SU(3)_C$ gauge couplings. The numerical solution gives $k \sim 0.5$.

The ratio $(\Lambda_{Z'}/M)$ is obtained from the one loop renormalization group equations

$$\left(\frac{\Lambda_{Z'}}{M}\right) = exp\left(\frac{32\pi(1-\alpha_5(M))}{b_5}\right) \tag{15}$$

with $b_5 = \frac{1}{2}n_5 - 15$, $\alpha_5(M) \sim 0.059$ and $n_5 = 8$ which give $\Lambda_{Z'} \sim 10^{14} GeV$. Taking the $SO(10)$ singlet VEVs to be $\sim 10^{-1}M$, a numerical scenario for the heavy generation gives

$$m_{\nu_\tau} \sim 25\ eV \qquad m_{N_\tau} \sim 100\ TeV \tag{16}$$

Assuming that other singlet masses arise at order $N = 10$ (or $m = 4$), one gets

the loose bounds for the other two generations:

$$m_{\nu_e}, m_{\nu_\mu} < O(10^{-4}\ eV) \quad m_{N_e}, m_{N_\mu} > O(1\ TeV) \tag{17}$$

If some of the $5, \bar{5}$ states of the hidden $SU(5)_H$ get mass after the symmetry breakdown, n_5 gets smaller. For smaller n_5, $\Lambda_{Z'}$ and the see-saw scale m_χ increase. Therefore the left-handed neutrino masses decrease and the right–handed neutrino masses increase. It is also possible to modify the boundary conditions of the hidden sector in the vector γ: $\gamma(\phi^{1,2\cdots 8}) = (\frac{1}{2}000\frac{1}{2}\frac{1}{2}\frac{1}{2}0)$ [15]. Then the hidden gauge group is $SU(7) \times U(1)^2$ in which case $\Lambda_{Z'}$ is larger and the left–handed neutrino masses are smaller.

4. Conclusions

We have shown that although there is a large Dirac mass for the $\tau-$ neutrino at tree level in the superpotential of the superstring standard–like models, very light left–handed and heavy right–handed neutrinos for all generations are obtained. This is due to the see–saw mechanism present in the model as a consequence of the higher order nonrenormalizable terms and the breaking of $U(1)_{Z'}$ symmetry by hidden sector condensates. The hidden sector condensates break $U(1)_{Z'}$ and leave $U(1)_Y$ unbroken. We learn from our superstring derived see–saw mechanism a new feature unprecedented in traditional GUT models. Although the scale of $U(1)_{Z'}$ breaking can be large (e.g. $\Lambda_{Z'} \approx 10^{14} GeV$) the effective see-saw scale is much smaller. This feature arises in superstring models because the see–saw and the singlet mass terms are obtained from higher order nonrenormalizable terms. As a result, the right–handed neutrino masses are suppressed relative to the $U(1)_{Z'}$ breaking scale. Consequently, the right–handed neutrinos can be at a rather low energy scale, say a few TeV, but this does not imply the existence of a new gauge boson (Z') at the TeV scale [16].

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	ψ^μ	$\{\chi^{12};\chi^{34};\chi^{56}\}$	$\bar\psi^1, \bar\psi^2, \bar\psi^3, \bar\psi^4, \bar\psi^5, \bar\eta^1, \bar\eta^2, \bar\eta^3$	$\bar\phi^1, \bar\phi^2, \bar\phi^3, \bar\phi^4, \bar\phi^5, \bar\phi^6, \bar\phi^7, \bar\phi^8$
α	0	{0, 0, 0}	1, 1, 1, 0, 0, 0, 0, 0	1, 1, 1, 1, 0, 0, 0, 0
β	0	{0, 0, 0}	1, 1, 1, 0, 0, 0, 0, 0	1, 1, 1, 1, 0, 0, 0, 0
γ	0	{0, 0, 0}	$\frac{1}{2}, \frac{1}{2}, \frac{1}{2}, \frac{1}{2}, \frac{1}{2}, \frac{1}{2}, \frac{1}{2}, \frac{1}{2}$	$\frac{1}{2}$, 0, 1, 1, $\frac{1}{2}, \frac{1}{2}, \frac{1}{2}$, 0

	$y^3y^6, y^4\bar y^4, y^5\bar y^5, \bar y^3\bar y^6$	$y^1\omega^6, y^2\bar y^2, \omega^5\bar\omega^5, \bar y^1\bar\omega^6$	$\omega^1\omega^3, \omega^2\bar\omega^2, \omega^4\bar\omega^4, \bar\omega^1\bar\omega^3$
α	1, 0, 0, 0	0, 0, 1, 1	0, 0, 1, 1
β	0, 0, 1, 1	1, 0, 0, 0	0, 1, 0, 1
γ	0, 1, 0, 1	0, 1, 0, 1	1, 0, 0, 0

Table 1. A three generations $SU(3)\times SU(2)\times U(1)^2$ model. The choice of generalized GSO coefficients is: $c\begin{pmatrix} b_j \\ \alpha,\beta,\gamma\end{pmatrix} = -c\begin{pmatrix}\alpha\\1\end{pmatrix} = c\begin{pmatrix}\alpha\\\beta\end{pmatrix} = -c\begin{pmatrix}\beta\\1\end{pmatrix} = c\begin{pmatrix}\gamma\\1,\alpha\end{pmatrix} = -c\begin{pmatrix}\gamma\\\beta\end{pmatrix} = -1$ (j=1,2,3), with the others specified by modular invariance and space–time supersymmetry. Trilevel Yukawa couplings are obtained only for $+\frac{2}{3}$ charged quarks. The 16 right–moving internal fermionic states $\{\bar\psi^{1,\cdots,5}, \bar\eta^1, \bar\eta^2, \bar\eta^3, \bar\phi^{1,\cdots,8}\}$, correspond to the 16 dimensional compactified torus of the ten dimensional heterotic string. The 12 left–moving and 12 right–moving real internal fermionic states correspond to the six left and six right compactified dimensions in the bosonic language. ψ^μ are the two space–time external fermions in the light–cone gauge and χ^{12}, χ^{34}, χ^{56} correspond to the spin connection in the bosonic constructions.

F	SEC	$SU(3)_C \times SU(2)_L$	Q_C	Q_L	Q_1	Q_2	Q_3	Q_4	Q_5	Q_6	$SU(5) \times SU(3)$	Q_7	Q_8
V_1	$b_1 + 2\gamma + (I)$	(1,1)	0	0	0	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	0	0	(1,3)	$-\frac{1}{2}$	$\frac{5}{2}$
$\bar{V}_1$		(1,1)	0	0	0	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	0	0	$(1,\bar{3})$	$\frac{1}{2}$	$-\frac{5}{2}$
T_1		(1,1)	0	0	0	$\frac{1}{2}$	$\frac{1}{2}$	$-\frac{1}{2}$	0	0	(5,1)	$-\frac{1}{2}$	$-\frac{3}{2}$
$\bar{T}_1$		(1,1)	0	0	0	$\frac{1}{2}$	$\frac{1}{2}$	$-\frac{1}{2}$	0	0	$(\bar{5},1)$	$\frac{1}{2}$	$\frac{3}{2}$
V_2	$b_2 + 2\gamma + (I)$	(1,1)	0	0	$\frac{1}{2}$	0	$\frac{1}{2}$	0	$\frac{1}{2}$	0	(1,3)	$-\frac{1}{2}$	$\frac{5}{2}$
$\bar{V}_2$		(1,1)	0	0	$\frac{1}{2}$	0	$\frac{1}{2}$	0	$\frac{1}{2}$	0	$(1,\bar{3})$	$\frac{1}{2}$	$-\frac{5}{2}$
T_2		(1,1)	0	0	$\frac{1}{2}$	0	$\frac{1}{2}$	0	$-\frac{1}{2}$	0	(5,1)	$-\frac{1}{2}$	$-\frac{3}{2}$
$\bar{T}_2$		(1,1)	0	0	$\frac{1}{2}$	0	$\frac{1}{2}$	0	$-\frac{1}{2}$	0	$(\bar{5},1)$	$\frac{1}{2}$	$\frac{3}{2}$
V_3	$b_3 + 2\gamma + (I)$	(1,1)	0	0	$\frac{1}{2}$	$\frac{1}{2}$	0	0	0	$\frac{1}{2}$	(1,3)	$-\frac{1}{2}$	$\frac{5}{2}$
$\bar{V}_3$		(1,1)	0	0	$\frac{1}{2}$	$\frac{1}{2}$	0	0	0	$\frac{1}{2}$	$(1,\bar{3})$	$\frac{1}{2}$	$-\frac{5}{2}$
T_3		(1,1)	0	0	$\frac{1}{2}$	$\frac{1}{2}$	0	0	0	$-\frac{1}{2}$	(5,1)	$-\frac{1}{2}$	$-\frac{3}{2}$
$\bar{T}_3$		(1,1)	0	0	$\frac{1}{2}$	$\frac{1}{2}$	0	0	0	$-\frac{1}{2}$	$(\bar{5},1)$	$\frac{1}{2}$	$\frac{3}{2}$

Table 2. Massless states from the sectors $b_j + 2\gamma$, and their quantum numbers.

F	SEC	$SU(3)_C \times SU(2)_L$	Q_C	Q_L	Q_1	Q_2	Q_3	Q_4	Q_5	Q_6	$SU(5) \times SU(3)$	Q_7	Q_8
H_{13}	$b_1 + b_3 + \alpha$	(1,1)	$-\frac{3}{4}$	$\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{4}$	0	0	0	(1,3)	$\frac{3}{4}$	$\frac{5}{4}$
H_{14}	$\pm\gamma + (I)$	(1,1)	$\frac{3}{4}$	$-\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$\frac{1}{4}$	0	0	0	$(1,\bar{3})$	$-\frac{3}{4}$	$-\frac{5}{4}$
H_{15}		(1,2)	$-\frac{3}{4}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{4}$	0	0	0	(1,1)	$-\frac{1}{4}$	$-\frac{15}{4}$
H_{16}		(1,2)	$\frac{3}{4}$	$\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$\frac{1}{4}$	0	0	0	(1,1)	$\frac{1}{4}$	$\frac{15}{4}$
H_{17}		(1,1)	$-\frac{3}{4}$	$\frac{1}{2}$	$-\frac{1}{4}$	$-\frac{3}{4}$	$-\frac{1}{4}$	0	0	0	(1,1)	$-\frac{1}{4}$	$-\frac{15}{4}$
H_{18}		(1,1)	$\frac{3}{4}$	$-\frac{1}{2}$	$\frac{1}{4}$	$\frac{3}{4}$	$\frac{1}{4}$	0	0	0	(1,1)	$\frac{1}{4}$	$\frac{15}{4}$
H_{19}	$b_2 + b_3 + \alpha$	(1,1)	$-\frac{3}{4}$	$\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	0	0	0	(5,1)	$-\frac{1}{4}$	$\frac{9}{4}$
H_{20}	$\pm\gamma + (I)$	(1,1)	$\frac{3}{4}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{4}$	0	0	0	$(\bar{5},1)$	$\frac{1}{4}$	$-\frac{9}{4}$
H_{21}		(3,1)	$\frac{1}{4}$	$\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	0	0	0	(1,1)	$-\frac{1}{4}$	$-\frac{15}{4}$
H_{22}		$(\bar{3},1)$	$-\frac{1}{4}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{4}$	0	0	0	(1,1)	$\frac{1}{4}$	$\frac{15}{4}$
H_{23}		(1,1)	$-\frac{3}{4}$	$\frac{1}{2}$	$\frac{1}{4}$	$-\frac{1}{4}$	$\frac{3}{4}$	0	0	0	(1,1)	$\frac{1}{4}$	$\frac{15}{4}$
H_{24}		(1,1)	$\frac{3}{4}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$\frac{1}{4}$	$-\frac{3}{4}$	0	0	0	(1,1)	$\frac{1}{4}$	$\frac{15}{4}$
H_{25}		(1,1)	$-\frac{3}{4}$	$\frac{1}{2}$	$\frac{1}{4}$	$\frac{3}{4}$	$-\frac{1}{4}$	0	0	0	(1,1)	$-\frac{1}{4}$	$-\frac{15}{4}$
H_{26}		(1,1)	$\frac{3}{4}$	$-\frac{1}{2}$	$-\frac{1}{4}$	$-\frac{3}{4}$	$\frac{1}{4}$	0	0	0	(1,1)	$\frac{1}{4}$	$\frac{15}{4}$

Table 3. **Massless scalars and their quantum numbers**