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**PROBLEMS FOR SUPERSTRING MODELS
WITH VACUUM EXPECTATION VALUES FOR CONJUGATE SNEUTRINOS**

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ABSTRACT

Superstring models compactified on Calabi-Yau manifolds contain two new neutral scalars in each 27 generation of chiral supermultiplets. We consider implications of a vacuum expectation value for one of the neutral chiral scalars, the conjugate sneutrino $\tilde{\nu}_L^c$: $\langle 0 | \tilde{\nu}_L^c | 0 \rangle \equiv y$. Apart from a coupling which generates a Dirac neutrino mass and which must be small [$\leq O(10^{-9})$], its only superpotential interaction is with $D_L d_L$, so there would be D-d mixing if $y \neq 0$. The fact that m_d is not much less than m_e or m_u then tells us that $y \leq O(\text{few TeV})$. When the four-dimensional gauge group has rank 5, a Pseudo-Goldstone boson a_ν arises if $y \neq 0$, and the D-d mixing would then induce the decay $K \rightarrow \pi + a_\nu$ at an unacceptable rate. Regardless of the gauge group rank, we show that the combined constraints on D-d mixing, and on the flavour-changing neutral currents it would induce, are sufficient to rule out $D_L d_L^c \tilde{\nu}_L^c$ superpotential coupling large enough to produce $\langle 0 | \tilde{\nu}_L^c | 0 \rangle \neq 0$. These problems are avoided in models with a rank 5 gauge group and $\langle 0 | \tilde{\nu}_L^c | 0 \rangle \leq 40 \text{ keV}$.

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With the advent of superstring theories, and in particular the $E_8 \otimes E_8$ heterotic superstring [1], as candidate theories of everything (TOEs), much activity has centred on the compactification of these theories to four-dimensional models of particle interactions [2]. A guiding principle generally imposed [3] on the compactification is that the resulting four-dimensional model maintains four-dimensional $N = 1$ supersymmetry, which will be used to preserve the hierarchy between the natural string and compactification scales around the Planck mass, and the low-energy scale $[O(100) \text{ GeV}]$ of the observed gauge interactions. Compactifications of the 10-dimensional superstring theory space-time to $M_4 \otimes K$, where M_4 is a maximally symmetric space-time and K is the 'internal' space, which have four-dimensional $N = 1$ supersymmetry necessarily use spaces K of $SU(3)$ holonomy [3]. If it is further assumed that there is no torsion in the spin connection, then such a manifold K must be Kähler with zero Ricci tensor: a Calabi-Yau manifold. If for the $E_8 \otimes E_8$ heterotic string [1] the observable sector gauge connection one form is set equal to that of the spin connection, in order to avoid gauge and gravitational anomalies and satisfy the string theory equations of motion [3]*, then the observable sector E_8 gauge group will be broken to E_6 by the compactification. If K is not simply connected, then the E_6 may be further broken by gauge fields winding around the 'holes' in K [6]. This can lead to non-zero values of Wilson line integrals for loops in non-trivial homotopy classes, which act like symmetry-breaking adjoint fields. The resulting low-energy gauge theory will have rank 6 if $\pi_1(K)$ is Abelian, but may have rank 5 for non-Abelian $\pi_1(K)$ [3,7]. It is possible, in principle, to obtain four-dimensional gauge groups of rank 4 or 5 by compactifying on manifolds in which the spin and gauge connections are not identified [8]. However, these manifolds have theoretical [9] and phenomenological [10] problems, and will not be discussed further here.

The chiral matter multiplets in the observable sector arise from zero modes of the E_8 gauge supermultiplet when compactified on K . In general there will be some number of generations N_g of 27-dimensional representations of the E_6 broken by Wilson lines, plus possibly some number δ of paired subsets of states from the $27 + \overline{27}$ representations, which depend on the details of the Wilson line symmetry-breaking, and form representations of the low-energy group [7]. Each 27 generation of chiral multiplets contains the following set of states, labelled by their $SU(3) \times SU(2) \times U(1)_Y \times U(1)_{Y'}$ quantum numbers [Y' is the linear combination of $Y_L + Y_R$, in the $SU(3)_c \times SU(3)_L \times SU(3)_R$ decomposition of E_6 , which is orthogonal to hypercharge Y]:

$$Q(3, 2, \frac{1}{3}, \frac{1}{3}), d^c(\overline{3}, 1, \frac{2}{3} - \frac{1}{6}), u^c(\overline{3}, 1 - \frac{4}{3}, \frac{1}{3}), L(1, 2, -1, -\frac{1}{6}),$$

$$e^c(1, 1, 2, \frac{1}{3}), H(1, 2, 1, -\frac{2}{3}), \overline{H}(1, 2, -1, -\frac{1}{6}),$$

$$D(3, 1, -\frac{2}{3}, -\frac{2}{3}), D^c(\overline{3}, 1, \frac{2}{3} - \frac{1}{6}), \nu^c(1, 1, 0, \frac{5}{6}), N(1, 1, 0, \frac{5}{6}).$$

(We take our chiral supermultiplets to be left-handed, and in the cases where one of the numbers of the supermultiplet is a Standard Model field we use that label for the supermultiplet.) As well as the 15 states of the Standard Model generation, we note the presence of two Higgs doublets ($H, \overline{H}$) of opposite isospin and $U(1)_Y$, an isosinglet charge $-\frac{1}{3}$ quark (D, D^c), and two $SU(3) \otimes SU(2) \otimes U(1)$ singlets (ν^c, N), in each 27 generation.

In this paper we argue that most models with $\langle 0|\tilde{\nu}^c|0\rangle \neq 0$ have phenomenological problems, which are avoided by rank 5 models with $\langle 0|\tilde{\nu}^c|0\rangle \lesssim 40 \text{ keV}$.

*) Even if Calabi-Yau manifolds are not exact solutions to the equations of motion [4], solutions with identical topological properties can be constructed from them [5].

The superpotential interaction among the chiral supermultiplets is a four-dimensional result of the gauge interactions of the E_6 gauge supermultiplet zero modes on K . As such, all renormalizable trilinear terms in the superpotential must be of a form that would appear in the E_6 invariant trilinear 27^3 , although the coefficients of the various terms in the expansion of 27^3 in multiplets of the low-energy gauge group in general only bear the relations imposed by the low-energy group. The most general set of trilinear superpotential terms involving fields of the 27 generations can be written [3, 7, 11]:

$$\begin{aligned}
W = & \lambda_0 H L \nu^c + \lambda_1 \bar{H} L e^c + \lambda_2 \bar{H} Q d^c + \lambda_3 H Q u^c \\
& + \lambda_4 N \bar{H} H + \lambda_5 N D^c D + \lambda_6 D Q Q + \lambda_7 D^c u^c d^c \\
& + \lambda_8 D^c Q L + \lambda_9 D u^c e^c + \lambda_{10} D d^c \nu^c
\end{aligned} \tag{1}$$

where each coupling λ is actually a tensor with three generation indices^{*)}. We will generally assume that these coupling tensors do not have conspiratorial zeros. Some of our arguments could be evaded by postulating convenient zeros for some particular values of these generation indices: we are not enamoured of such a conspiracy theory. Constraints on the superpotential from baryon and lepton number conservation, neutrino masses, and neutral-flavour conservation are analysed elsewhere [11]; for our purposes it suffices to examine the superpotential couplings of the $SU(3) \otimes SU(2) \otimes U(1)$ invariant supermultiplets:

$$W_N = \lambda_4 N \bar{H} H + \lambda_5 N D^c D \tag{2a}$$

$$W_{\nu^c} = \lambda_0 H L \nu^c + \lambda_{10} D d^c \nu^c \tag{2b}$$

Vacuum expectation values for the neutral scalars ($\langle 0|N|0 \rangle \equiv x$, $\langle 0|\tilde{\nu}^c|0 \rangle \equiv y$) would manifest themselves through these superpotential terms, as well as through the gauge symmetry breaking they would induce. There are at least 3 N and 3 $\tilde{\nu}^c$ fields in the effective low-energy theory, and perhaps more if one invokes split multiplets. We can always work in a field basis [12] where at most one representative of the fields with a given set of quantum numbers has a non-vanishing v.e.v.

Since the Standard Model has a rank 4 gauge group, then for models with rank 5 four-dimensional gauge group at least one of x, y must be non-zero, and for models with a rank 6 gauge group they both must be non-vanishing. In general we [13] do expect $\langle 0|N|0 \rangle = x \neq 0$; as well as reducing the rank of the unbroken gauge group by one, its superpotential couplings are responsible for Higgs mixing ($\lambda_4 N \bar{H} H$), and for supplying a mass to the isoscalar D -quarks, $\lambda_5 N D^c D$. Furthermore, these couplings induce terms in the renormalization group equation for the N mass squared, which can drive it negative if they are sufficiently strong, and so trigger low-energy spontaneous symmetry-breaking [13]. Note that the N v.e.v. x should not be much larger than

^{*)} If there are light 'survivors' from split $27 + \bar{27}$ representations, the indices on the couplings λ should be increased to include light fields from additional 27 's.

$O(1 \text{ TeV})$, as that is the scale of the (necessary) Higgs mass mixing it induces. This also means that $m_D = O(1 \text{ TeV})$. So the options which we must consider are: i) the case where only the N scalar has a non-zero v.e.v. — and the gauge group is of necessity rank 5, to avoid an unobserved, unbroken $U(1)$ gauge group; and ii) the cases where both the N and $\bar{\nu}$ have non-zero v.e.v.'s with a gauge group of rank either 5 or 6. It is important to note that since we have required that our superpotential terms be those which would appear in the E_6 invariant 27^3 , the resulting superpotential will certainly be invariant under transformations of the rank 6 Cartan subalgebra of E_6 . If only five of those invariances are gauged, as in a rank 5 low-energy gauge group, then there will be a global, non-anomalous $U(1)$ symmetry of the superpotential left ungauged, whose consequences we now investigate.

There are only two superpotential terms involving the ν^c in Eq. (26). The first of these, $\lambda_0 H L \nu^c$, will give rise to fermionic bilinears $\lambda_0 H (L \nu^c)$ and $\lambda_0 \bar{\nu}^c (\bar{H} L)$, where $(\)$ denotes the Lorentz scalar bilinear combination of the fermion fields enclosed. In the presence of a Higgs v.e.v. $\langle 0|H|0 \rangle = v = O(100 \text{ GeV})$ the first term produces a Dirac neutrino mass $m_\nu = \lambda_0 v$, and with a sneutrino v.e.v. the second term produces a lepton-Higgs mixing mass term $m_{HL} = \lambda_0 y$. Since there is no ν^c bilinear coupling in the renormalizable superpotential Eq. (1) that could serve as a source of Majorana mass for the right-handed neutrino, there is no 'see-saw' mechanism present, and the Dirac masses produced by the $\lambda_0 H L \nu^c$ coupling will be the physical masses of the light neutrinos. Non-renormalizable terms might also appear in the superpotential, but these appear unable to produce a large Majorana mass [14] in the class of manifolds considered here. Limits on the neutrino masses from particle physics and astrophysics demand that they be less than $O(100 \text{ eV})$ [15]; this bounds $\lambda_0 = m_\nu/v \lesssim 10^{-9}$. Moreover, bounds on the lepton number violation induced by H-L mixing require $\lambda_0 \lesssim 10^{-2} (10^2 \text{ GeV}/y)$ [16]. So this coupling is almost if not entirely absent from the superpotential, and is in any case too small to drive $m_\nu^2 < 0$ and thereby generate $\langle 0|\nu^c|0 \rangle = y \neq 0$.

The other possible superpotential coupling of ν^c is $\lambda_{10} D d^c \nu^c$. This term is phenomenologically acceptable provided that it does not appear in conjunction with $\lambda_6 D Q Q + \lambda_7 D^c u^c d^c$ terms—which would induce proton decay [7,11]. In combination with $\lambda_8 D^c Q L + \lambda_9 D u^c e^c$ terms it would induce neutrino masses by radiative corrections which might be acceptably small [11]*. With such a coupling, a $\bar{\nu}^c$ v.e.v. would give rise to a mass mixing between the d and D quarks $\lambda_{10} y (D d^c)$. The full charge $-1/3$ quark mass matrix would then be (suppressing generation indices):

$$(d^c, D^c) \begin{pmatrix} \lambda_2 \bar{\nu} & \lambda_{10} y \\ 0 & \lambda_5 x \end{pmatrix} \begin{pmatrix} d \\ D \end{pmatrix} + h.c. \quad (3)$$

which may be diagonalized by a bi-unitary transformation of the chiral fermion fields [11]:

$$\begin{pmatrix} m_< & 0 \\ 0 & m_> \end{pmatrix} = \begin{pmatrix} C_R & S_R \\ -S_R & C_R \end{pmatrix} \begin{pmatrix} \lambda_2 \bar{\nu} & \lambda_{10} y \\ 0 & \lambda_5 x \end{pmatrix} \begin{pmatrix} C_L & S_L \\ -S_L & C_L \end{pmatrix} \quad (4)$$

*) Either of these generation-independent patterns may arise for topological reasons, or if the compactification exhibits a discrete symmetry as simple as Z_2 [11].

where $c_{L,R} = \cos(\theta_{L,R})$, $s_{L,R} = \sin(\theta_{L,R})$, and we suppress the generation structure of the mass matrix. The mass eigenvalues are

$$m_{<} = c_R (\lambda_2 \bar{v} c_L - \lambda_{10} y s_L) - s_R \lambda_5 x s_L \quad (5a)$$

$$m_{>} = -s_R (\lambda_2 \bar{v} s_L + \lambda_{10} y c_L) + c_R \lambda_5 x c_L \quad (5b)$$

with the mixing angles determined by [11]

$$\tan \theta_L \simeq y \lambda_2 \lambda_{10} \bar{v} / (\lambda_5^2 x^2 + \lambda_{10}^2 y^2) \quad (6a)$$

$$\tan \theta_R \simeq -\lambda_{10} y / \lambda_5 x \quad (6b)$$

where we have assumed $\lambda_2 \bar{v} \ll (\lambda_5^2 x^2 + \lambda_{10}^2 y^2)^{1/2}$ as required by phenomenology: $\theta_L \ll 1$ and $m_{>}^2 > m_{<}^2$.

The approximate mass eigenvalues are then

$$m_d \equiv m_{<} \simeq \lambda_2 \bar{v} \lambda_5 x / (\lambda_5^2 x^2 + \lambda_{10}^2 y^2)^{1/2} \quad (7a)$$

$$m_D \equiv m_{>} \simeq (\lambda_5^2 x^2 + \lambda_{10}^2 y^2)^{1/2} \quad (7b)$$

where we identify the known charge $-1/3$ quarks with the lighter states, so $m_{<} \sim 1$ GeV. *A priori* their right-hand components could be either mostly d^c ($|\sin \theta_R| < 1/\sqrt{2}$) or mostly D^c ($|\sin \bar{\theta}_R| \equiv |\sin(\pi/2 - \theta_R)| < 1/\sqrt{2}$). The charged leptons and charge $2/3$ quarks also obtain masses through Yukawa couplings to Higgs v.e.v.'s. In the lepton case it is the same $\bar{H}$ v.e.v., and the mass obtained is $\lambda_1 \bar{v}$ which is again of order 1 GeV. We expect the superpotential couplings of leptons, and of charge $-1/3$ quarks, to $\bar{H}$ to be comparable at the unification scale, and hence, after renormalization group evolution to low energies, to agree within an order of magnitude. Therefore we certainly expect $\lambda_2 \bar{v} \lesssim 10(\lambda_1 \bar{v})$. Then, since physical charge $-1/3$ quarks are certainly no lighter than charged leptons,

$$|m_{<}| \simeq \frac{\lambda_2 \bar{v} \lambda_5 x}{(\lambda_5^2 x^2 + \lambda_{10}^2 y^2)^{1/2}} \gtrsim m_L \quad (8)$$

we find from equation (7a) that

$$\lambda_{10} y \lesssim 10 (\lambda_5 x) \quad (9)$$

or for equal superpotential couplings $y \lesssim 10 x \approx O(\text{few TeV})$. Thus, even if $y \neq 0$, it cannot be very large.

The mass mixing generated by $y \neq 0$ will induce flavour-changing neutral current couplings of the Z^0 boson of the form $\sqrt{g_2^2 + g'^2} \theta_L^2 \bar{s}_L \gamma^\mu d_L Z_\mu$ [11], since D_L and d_L have different weak isospin and hypercharge. There will also be a similar coupling to the Z' boson (or bosons). As long as $x \lesssim 1 \text{ TeV}$ their will be at least one Z' with a mass in the few hundred GeV range that couples to a $U(1)$ current for which D_R and d_R have different charges, so that d_R/D_R mixing will induce couplings of the form $g' \sin^2 \theta_M \bar{s}_R \gamma^\mu d_R Z'_\mu$, where θ_M is the smaller of θ_R and $\bar{\theta}_R$. Detailed analysis [11] of the effects of these flavour-changing neutral currents on the processes $K^0 \leftrightarrow \bar{K}^0$, $B^0 \leftrightarrow \bar{B}^0$, $K \rightarrow \pi \nu \bar{\nu}$, and $K_L \rightarrow \bar{\mu} \mu$ shows that the experimental limits on these processes require: $\theta_L < 10^{-2}$ ($K^0 \leftrightarrow \bar{K}^0$, $B^0 \leftrightarrow \bar{B}^0$, $K \rightarrow \pi \nu \bar{\nu}$, $K_L \rightarrow \bar{\mu} \mu$), and $\theta_M < 3 \times 10^{-2}$ ($K \rightarrow \pi \nu \bar{\nu}$, $K_L \rightarrow \bar{\mu} \mu$) or $\theta_M < 2 \times 10^{-2}$ ($K^0 \leftrightarrow \bar{K}^0$, $B^0 \leftrightarrow \bar{B}^0$). But for small $\bar{\theta}_R \approx \lambda_5 x / \lambda_{10} y$, $\bar{\theta}_R < 2 \times 10^{-2}$ demands $\lambda_{10} y > 50 \lambda_5 x$, in direct contradiction with the constraint (9). We are thus led to assume a small D_R component of the right-handed charge $-1/3$ quarks, and $\theta_R < 2 \times 10^{-2}$ requires $\lambda_{10} y < 1/50 \lambda_5 x$.

In the case of a rank-5 low-energy gauge group, the renormalizable terms in the superpotential (1) are invariant under a global $U(1)$ corresponding to the ungauged generator in the Cartan subalgebra of E_6 . If there were no other terms in the effective low-energy theory, we would have a Goldstone boson a_ν if both x and y were non-zero. In fact, we expect the low-energy theory to contain non-renormalizable terms of higher mass dimension in the superpotential, scaled by the appropriate inverse powers of m_P . These terms should respect the residual rank-5 low-energy gauge invariance, but need not be invariant under the global $U(1)$ corresponding to the ungauged sixth generator of the Cartan subalgebra. Therefore a_ν is in fact a pseudo-Goldstone boson. It is easy to check that the first possible superpotential terms of this type would have mass dimension 6 scaled by m_P^{-3} . These can give a mass to the a_ν at most of order $m_W^{5/2} / m_P^{3/2} < 10^{-20} \text{ GeV}$. For phenomenological purposes, therefore, the a_ν can be considered as essentially massless.

The ungauged global $U(1)$ which is not broken by v , $\bar{v} \neq 0$ is in fact $B - L$. In this case it is also unbroken by $x \neq 0$, and the Goldstone boson a_ν generated by $\langle 0 | \bar{\nu}^c | 0 \rangle \neq 0$ couples to the divergence of the $B - L$ current. If there is d/D mixing this current acquires $O(\theta_L^2, \theta_R^2)$ mixing among charged $-1/3$ quarks, and an effective flavour changing interaction is induced:

$$\mathcal{L}_{K\pi} = F_{K\pi}^{-1} K \sum_\mu \pi \partial^\mu a \quad (10)$$

For an effective Lagrangian of this form, Wilczek [17] obtains

$$\frac{\Gamma(K^+ \rightarrow \pi^+ a)}{\Gamma(K^+ \rightarrow \pi^+ \pi^0)} = \frac{1.3 \times 10^{14} \text{ GeV}^2}{F_{K\pi}^2} \quad (11)$$

and using the experimental limit [18] $BR(K^+ \rightarrow \pi^+ a) < 3.8 \times 10^{-8}$ we have $F_{K\pi} > 2.7 \times 10^{10} \text{ GeV}$. For the phenomenologically acceptable case $\theta_R \ll 1$ we have

$$F_{K\pi}^{-1} = (\sin^2\theta_L + \sin^2\theta_R)/y \approx \sin^2\theta_R/y \approx \lambda_0^2 y / \lambda_5^2 x^2 \quad (12)$$

which, for $\lambda_{10} \sim \lambda_5$ and $x \leq 1$ TeV, requires from the bound on $F_{K\pi}$: $y \leq 40$ keV. One can of course allow y large with very small λ_{10} , but as will be seen below there is no mechanism for $y \neq 0$ if λ_{10} is very small.

We will now argue that regardless of the gauge group rank, one cannot generate $\langle 0|\tilde{\nu}^c|0\rangle = y \neq 0$ in dynamical models based on the superstring. Since in models with rank 6 gauge groups one needs $y \neq 0$ in order to give mass to the extra low energy gauge bosons, this poses a severe difficulty for models of this type. The first step in our argument is to note that there are no terms in the potential linear in $\tilde{\nu}^c$. Such terms could in principle arise from soft supersymmetry breaking which [19], given a superpotential, induces terms in the effective scalar potential which are functions of the scalar components alone that take the same form. The potential terms which could be so induced, $\tilde{D}\tilde{d}^c\tilde{\nu}^c$ or $H\tilde{L}\tilde{\nu}^c$, involve fields with no v.e.v. and hence will not produce a potential linear in $\tilde{\nu}^c$. In the absence of potential terms linear in $\tilde{\nu}^c$, we require that $m_{\tilde{\nu}^c}^2 < 0$ in order to produce a non-zero $\tilde{\nu}^c$ v.e.v. A negative mass squared for $\tilde{\nu}^c$ is required to drive $y \neq 0$, both in models where dimensional transmutation generates a small value of $\langle 0|\tilde{\nu}^c|0\rangle$ and also in those where the $\tilde{\nu}^c$ v.e.v. develops along D-flat (and F-flat) directions of the low-energy theory, and is driven to an intermediate scale. In order to drive $m_{\tilde{\nu}^c}^2 < 0$ via renormalization group scaling, we require Yukawa couplings in the renormalization group equation [13, 20], of strength roughly comparable to the $U(1)'$ gauge interactions, which have the opposite effect. Yukawa couplings of this strength for the $\tilde{\nu}^c$ cannot arise from the $\lambda_0 H L \nu^c$ term: we have already seen that limits on neutrino masses demand $\lambda_0 \leq 10^{-9}$. The other possibility is the $D d^c \nu^c$ coupling. We have already seen that the requirement of reasonable masses for charge $-1/3$ quarks together with constraints from induced flavour-changing neutral current processes constrains $\lambda_{10} y < 1/50 \lambda_5 x$. However if $B - L$ is gauged, the associated Z' is appreciably more massive than the Z : $y \geq 1/2$ TeV. Then for $x \leq 1$ TeV we must require $\lambda_{10}/\lambda_5 \leq 1/25$. Since no Yukawa couplings are expected to be strong, λ_{10} is bounded to be very weak and cannot drive $m_{\tilde{\nu}^c}^2$ negative.

We have presented arguments against a vacuum expectation value for the $\tilde{\nu}^c$ field in Calabi-Yau compactifications of superstring theories. As we have already mentioned, it is possible that one could circumvent these arguments by postulating convenient zeros in the Yukawa coupling matrices, which would follow from discrete symmetries. If one does not try to follow this route, the absence of such a v.e.v. implies strong constraints on the models one may construct. First, the absence of such a v.e.v. means that gauge symmetry breaking outside the Standard Model may only occur via N v.e.v.'s. This indicates that the gauge group produced by Wilson line breaking of E_6 cannot have rank larger than 5. In a previous paper [10] some of us have argued that models with no gauge interactions for the N field, including *a fortiori* models with rank 4, have severe problems. Only models of rank 5 in four dimensions survive the arguments there and in this paper. In the case of Calabi-Yau manifolds, there is a unique rank 5 group which one may obtain by Wilson line breaking of E_6 [7], so this fixes the gauge group. It may prove possible to arrive at other rank 5 groups using orbifold compactifications [21], but this remains to be seen.

Since the N v.e.v. must be of order one TeV so as to give the correct Higgs mixing, and since it is the only v.e.v. outside the Standard Model, we find that all symmetry breaking below the compactification scale occurs at a scale of $O(1 \text{ TeV})$ or less; there can be no intermediate scale.

In the absence of intermediate scales, to have viable renormalization group evolution of the gauge coupling constants [7], we are restricted to models with three generations.

Thus we see that the absence of a vacuum expectation value for the conjugate sneutrino, as argued in this paper, has important consequences.

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REFERENCES

- [1] D.J. Gross, J. Harvey, E. Martinec and R. Rohm, Phys. Rev. Lett. **54** (1985) 502; Nucl. Phys. **B256** (1985) 253 and **B267** (1986) 75.
- [2] J. Ellis, Proc. 7th Workshop on Grand Unification, ICOBAN '86, CERN preprint TH. 4439/86 (1986).
- [3] P. Candelas, G.T. Horowitz, A. Strominger and E. Witten, Nucl. Phys. **B258** (1985) 46.
- [4] M.T. Grisaru, A. van der Ven and D. Zanon, Harvard/Brandeis preprints HUTP-86/A020, A026, A027 (BRX-TH-196, 198, 199) (1986).
D.J. Gross and E. Witten, Princeton Univ. preprint (1986).
M.D. Freeman and C.N. Pope, Imperial College preprint TP/85-86/17 (1986).
- [5] D. Nemeschansky and A. Sen, SLAC preprint PUB-3925 (1986).
A. Sen, SLAC preprint PUB-3956 (1986).
- [6] Y. Hosotani, Phys. Lett. **129B** (1985) 193.
- [7] E. Witten, Nucl. Phys. **B258** (1985) 75.
S. Cecotti, J.-P. Derendinger, S. Ferrara, L. Girardello and M. Roncadelli, Phys. Lett. **156B** (1985) 318.
J.D. Breit, B.A. Ovrut and G. Segrè, Phys. Lett. **158B** (1985) 33.
A. Sen, Phys. Rev. Lett. **55** (1985) 33.
M. Dine, V. Kaplunovsky, M. Mangano, C. Nappi and N. Seiberg, Nucl. Phys. **B259** (1985) 519.
- [8] E. Witten, Nucl. Phys. **B268** (1986) 79.
- [9] J. Ellis, C. Gómez, D.V. Nanopoulos and M. Quirós, Phys. Lett. **173B** (1986) 59.
M. Dine, N. Seiberg, X.-G. Wen and E. Witten, Princeton preprint (1986).
- [10] J. Ellis, K. Enqvist, D.V. Nanopoulos, K. Olive, M. Quirós and F. Zwirner, CERN preprint TH. 4449/86 (1986).
- [11] B.A. Campbell, J. Ellis, K. Enqvist, M.K. Gaillard and D.V. Nanopoulos, CERN preprint TH. 4473/86 (1986).
- [12] H. Georgi and D.V. Nanopoulos, Phys. Lett. **82B** (1979) 95.
- [13] E. Cohen, J. Ellis, K. Enqvist and D.V. Nanopoulos, Phys. Lett. **161B** (1985) 85.
J. Ellis, K. Enqvist, D.V. Nanopoulos and F. Zwirner, Nucl. Phys. **B276** (1986) 14;
Mod. Phys. Lett. **A1** (1986) 57.
- [14] J.P. Derendinger, L. Ibáñez, and H.P. Nilles, Nucl. Phys. **B267** (1985) 365.
- [15] S. Sarkar and A.M. Cooper, Phys. Lett. **148B** (1984) 347.
- [16] L. Hall and M. Suzuki, Nucl. Phys. **B231** (1984) 419.
- [17] F. Wilczek, Phys. Rev. Lett. **49** (1982) 1549.
- [18] Y. Asano et al., Phys. Lett. **107B** (1981) 159.
- [19] R. Arnowitt, A.H. Chamseddine and P. Nath, Phys. Rev. Lett. **49** (1982) 970 and **50** (1983) 232;
Phys. Lett. **121 B** (1983) 33.
R. Barbieri, S. Ferrara and C. Savoy, Phys. Lett. **119B** (1982) 343.
H.-P. Nilles, M. Srednicki and D. Wyler, Phys. Lett. **120B** (1983) 345.
L.J. Hall, J. Lykken and S. Weinberg, Phys. Rev. **D27** (1983) 2359.
- [20] K. Inoue, A. Kakuto, H. Komatsu and S. Takeshita, Prog. Theor. Phys. **68** (1982) 927 and **71** (1984) 413.
- [21] L. Dixon, J.A. Harvey, C. Vafa and E. Witten, Nucl. Phys. **B261** (1985) 678;
Strings on orbifolds II, Princeton Univ. preprint (1986).