Intermediate scales of symmetry breaking in Calabi-Yau models*

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ABSTRACT

We discuss the generation of large intermediate scales of symmetry breaking in grand unification models suggested by the heterotic string. We analyze on dimensional grounds a particular scenario where one flat direction in the effective potential defines two different scales of gauge-symmetry breaking at very high energies [$\approx O(10^{15} \text{ GeV})$]. This mechanism implies the presence of one light [$O(1 \text{ TeV})$] nonchiral neutrino. The size of the observable low-energy effects seems, however, quite model dependent.

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MASTER

I. INTRODUCTION

The 10-dimensional $E_8 \times E_8$ heterotic string¹ contains near the Planck scale more dimensions, symmetry and matter than the observed at low energies. To obtain realistic models, it is necessary to assume that the (unknown) dynamics of the theory favours a vacuum that compactifies the six extra dimensions, breaks the excess of symmetry, and makes the extra fields heavy enough.

In the framework of compactification on a 3-generation Calabi-Yau manifold,² one may distinguish several processes that decouple from the model its exotic ingredients. The identification of the spin connection of the manifold in the gauge fields leaves in the observable sector $N = 1$ supersymmetry (SUSY) and E_6 gauge symmetry, with chiral superfields in the 27 and $\overline{27}$ representations of this group. In the known cases, the 3-family manifold is constructed from a non-simply connected one, where non-trivial gauge configurations provide a breaking of the remaining E_6 symmetry to a smaller rank-six group.³ This process will introduce a convenient asymmetry (not present in the E_6 model) between the quark and lepton sectors. The effective theory must also break SUSY. In the most promising scenario,⁴ SUSY is broken at large scales in a sector connected just through gravitational interactions with quarks and leptons and one obtains the usual (order 1 TeV and universal at the Planck scale) soft-breaking terms. Finally, in order to reduce the gauge symmetry to the rank-four group $SU(3)_C \times SU(2)_L \times U(1)_Y$, the evolution of the model down to low energies must define intermediate scales (IS) of symmetry breaking. Below these scales the exotic fields in the three chiral 27 of E_6 (two type *down* quarks, two neutrinos and two Higgs doublets) will combine into nonchiral representations of the Standard-Model symmetry and will (possibly) become massive. The vectorlike $27 + \overline{27}$ multiplets will also acquire masses through effective nonrenormalizable interactions. The IS must be very large (typically $10^{15} GeV$), since a slow proton decay and an acceptable value of the low-energy gauge couplings require that we are left at these energies with essentially a minimal supersymmetric extension of the Standard Model.⁷

Although the IS represent another source of uncertainty in the connexion of the string with its observable limit, it has been argued⁸ that they have model-independent implications that may be accesible to experiment. In particular, the IS would imply the existence of $O(1 TeV)$ fields producing neutrino masses and lepton-number violating processes. We would like here briefly review the mechanism that generates these IS and also introduce a particular structure in the superpotential (P) with some interesting properties. This structure requires just one flat direction to produce the two stages of symmetry breaking required in rank-six models. It provides differences of several orders of magnitude between the IS. In a realistic low-energy model, the Higgs sector (whose lightness is not protected by chirality⁷) and the neutrino sector (where see-saw masses⁵ would reflect the

physics in the *desert*) are in principle sensitive to the splitting of the scales. This particular structure is not contained in the analysis done by Arnowitt and Nath,⁸ where two independent flat directions are implicitly assumed. In contrast with them, we find only one nonchiral neutrino necessarily light [$O(1 \text{ TeV})$] and other quite model-dependent observable effects.

When we need to be definite we will refer to the Tian-Yau 3-generation manifold⁶⁻¹⁴ with

$$SU(3)_C \times SU(3)_L \times SU(3)_R \subset E_6 \quad (1.1)$$

gauge symmetry, which is the first and most extensively studied case. However, our arguments will also be relevant for the other 3-generation models (Schimmrigk and *bicubic* in $P^2 \times P^2$ ¹⁵). In the Tian-Yau model the non-singlet-gauge matter consists of nine families of leptons λ , six of $\bar{\lambda}$, seven of quarks q and antiquarks Q , and four of $\bar{q}$ and $\bar{Q}$, where

$$27 \longrightarrow \lambda \equiv (1, \bar{3}, 3) + q \equiv (3, 3, 1) + Q \equiv (\bar{3}, 1, \bar{3}) \quad (1.2)$$

and the assignment of standard quantum numbers is

$$\lambda \sim \begin{pmatrix} h^0 & h'^- & e \\ h^+ & h'^0 & \nu \\ e^c & \nu_4 & \nu_5 \end{pmatrix}; \quad q \sim \begin{pmatrix} u \\ d \\ d' \end{pmatrix} \times 3 \text{ colours}; \quad Q \sim \begin{pmatrix} u^c \\ d^c \\ d'^c \end{pmatrix} \times 3 \text{ colours}. \quad (1.3)$$

II. GENERATION OF INTERMEDIATE SCALES

Our starting point is the effective model at the compactification/unification scale M_c ($< M_{Planck}$). We suppose that the residual $N = 1$ SUSY is broken in the effective potential by scalar masses and by a term proportional to P . The SUSY-breaking effects would also fix the vacuum expectation values (VEV's) of the moduli fields (whose potential is in principle flat). These VEV's determine the complex structure of the manifold and then the value of all the couplings in P . We will assume that P incorporates a group of discrete symmetries acting in a definite way over each multiplet of chiral superfields. In Calabi-Yau models these symmetries are necessary to achieve the hierarchies M_Z/M_c or m_{up}/m_{top} ,⁹ and also to define a low-energy matter parity.¹⁰

Since the supersymmetric part of the effective potential is positive semidefinite, a nontrivial minimum should be favoured by the SUSY-breaking terms. Although these effective terms are in principle flavour independent (they are generated via gravitational interactions), their universality is broken by Yukawa radiative corrections. In Ref. 11 is showed via the renormalization group that the mass coefficients evolve fast (in one or two orders of magnitude) to negative values when one scales down the effective model, triggering nonzero VEVs. Then, if the potential contains a flat direction (a direction in the field space without self-interactions) the VEVs will grow up to values as large as the

scale where these mass coefficients become negative. It is remarkable that the Tian-Yau model mentioned above contains more families and couplings of leptons λ than of quarks q and Q , what favours VEV's for fields in lepton multiplets and preserves the colour symmetry. (A larger number of λ s above the IS is also necessary for a correct unification, since the gauge couplings verify $\alpha_C > \alpha_{L,R}$.)

The flatness of the potential is a necessary ingredient to define large IS. The supersymmetric scalar potential includes the D and F terms:

$$D^\alpha = g \phi_i^\dagger T^{\alpha i}_j \phi^j, \quad F_{\Phi_i} = \frac{\partial P}{\partial \Phi_i}, \quad (2.1)$$

$$V(\phi) = \sum_i |F_{\Phi_i}|^2 + \sum_\alpha |D^\alpha|^2. \quad (2.2)$$

D -flatness is obtained by requiring that for each VEV along a gauge flavour there is an identical VEV along a field in the conjugate representation; since these fields have opposite charges, their D -contributions cancel. F -flatness is obtained if all the terms in P contain more than one field with zero VEV. When only ν_5 and ν_4 [see (1.3)] develop VEV's, the gauge symmetry provides flatness with respect to trilinears in P . However, including effective $(27 \bar{27})^n$ nonrenormalizable terms (which we assume suppressed by inverse powers of M_c), the nonflat contributions can be forbidden only with the help of the discrete symmetries of the model.

In general, it will be difficult to find models with completely flat directions. Suppose that the VEV's along ν_5 involves n families; there are $n+1$ equations ($D=0$ and $F_i=0$, with $i=1, \dots, n$) but just n VEV's, and then a flat direction would require two dependent equations. The case with only one dependent equation (like in Ref. 12) corresponds to a discrete set of zeroes, with cancellations between terms of different dimension that would imply VEV's order M_c . Actually, completely flat directions have been found only in models that incorporate an R -symmetry.^{13,14} Experience shows, however, that when flatness is protected by an unbroken symmetry one also gets a large amount of extra massless fields, which tend to move the electroweak angle and the proton lifetime to non-acceptable values (this is the case, for instance, in Ref. 13).

Another possibility, easier to realize in models without R -symmetries, is the generation of IS along non-completely flat directions. If flatness is broken just by terms of high dimension (very suppressed by powers of M_c^{-1}), the scales may be large enough. Suppose that the effective potential receive contributions from terms $(27 \bar{27})^n$ in P . On dimensional grounds, the minimum will result from the balance between order $m_s^2 \approx 1 \text{ TeV}^2$ (negative) bilinears and order M_c^{-4n+6} (positive) F -contributions of dimension $4n-2$. The VEV's, order $^{2n-2} \sqrt{m_s/M_c} M_c$, will be one or two orders of magnitude below M_c (as required by perturbative unification of the gauge couplings) for $n=4-6$. (The soft

contribution proportional to P , order $m_s M_c^{2n-3}$, is unnecessary in this argument, because it is not the dominant for any value of the fields.)

We note that in the breaking along an *almost* flat direction the nonzero value of the F terms introduce a new source of SUSY breaking that will involve the scalars interacting with the fields developing VEV's. Contrarily to the ones coming from gravitational interactions, these soft-breaking terms are not universal and do not affect to the three chiral families (massless after the IS), although they could affect to the Higgs fields. We also observe that the fermionic partners of the scalars developing a VEV (non-weakly interacting neutrinos) receive only order m_s mass contributions from the terms $(27 \overline{27})^n$ in P . However, the Yukawa couplings of these neutrinos with the low-energy fields are necessarily very small (order $m_s/IS \approx 10^{-12}$); sizeable trilinears [of type $\nu_5(h^0 h^+)(h'^- h'^0)$] in P would give large masses to the low-energy fields.

Finally, we would like to mention the cosmological scenario that may accommodate this type of phase transitions.^{12,19} The finite-temperature effective potential would contain a quadratic term proportional to T^2 , where T is the temperature of the thermal bath in contact with the scalar field. For the field with flat potential that we are considering, this term prevents the transition (from zero to large VEV's) until the temperature of the Universe is $O(m_s \approx 1 \text{ TeV})$. Then, the entropy generated would dilute to non-acceptable values any existing baryon asymmetry. There are several ways to avoid this problem: to suppose a period of inflation (where the Universe is extremely cool before reheating and baryogenesis); to recreate the baryon number at $T = O(1 \text{ TeV})$ (transferring energy to heavy fields via anharmonic couplings); or simply to assume that the initial values of the field [at $T = O(M_c)$] are likely to be large (and the field would be trapped when T drops below the IS).

III. A MODEL WITH TWO INTERMEDIATE SCALES

The rank-six models under study here require two IS of gauge-symmetry breaking, defined by VEV's of two independent combinations of ν_4 and ν_5 (with identical $\overline{\nu}_4$ and $\overline{\nu}_5$ VEV's) in two different $\lambda + \overline{\lambda}$ vectorlike multiplets. It will be possible to make a gauge transformation and leave the multiplet containing bigger VEV's with $\langle {}^1\nu_4 \rangle = 0$ (the superindexes label the two λ and $\overline{\lambda}$). The smaller VEV's in the second family will in principle involve both gauge flavours. However, the fields ${}^2\nu_5$ and ${}^2\overline{\nu}_5$ are not neutral respect the possible matter parities one can define in the Tian-Yau manifold (see Table II). Since a low-energy matter parity is a necessary ingredient of any realistic SUSY model, these fields should not develop VEV's. We will consider models where this possibility is favoured by the discrete symmetries, and the VEV's in ${}^{1,2}\lambda$ and ${}^{1,2}\overline{\lambda}$ satisfy:

$$\langle {}^1\nu_5 \rangle = \langle {}^1\overline{\nu}_5 \rangle = \phi_1, \quad \langle {}^1\nu_4 \rangle = \langle {}^1\overline{\nu}_4 \rangle = 0; \quad (3.1a)$$

$$\langle {}^2\nu_5 \rangle = \langle {}^2\overline{\nu}_5 \rangle = 0, \quad \langle {}^2\nu_4 \rangle = \langle {}^2\overline{\nu}_4 \rangle = \phi_2, \quad (3.1b)$$

We define $x \equiv \phi_1/M_c$ and $y \equiv \phi_2/M_c$ ($1 > x > y$).

In the Tian-Yau model with $SU(3)_C \times SU(3)_L \times SU(3)_R$ gauge symmetry, the first VEV's [in Eq. (3.1a)] break the symmetry to $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, nine gauginos combine with the fermions $^1e, ^1\bar{e}, ^1e^c, ^1\bar{e}^c, ^1\nu, ^1\bar{\nu}, ^1\nu_4, ^1\bar{\nu}_4$, and $\frac{1}{\sqrt{2}}(^1\nu_5 - ^1\bar{\nu}_5)$ in the first supermultiplets, defining Dirac fields. At the second scale [ϕ_2 in (3.1b)] other three gauginos combine with $^2e^c, ^2\bar{e}^c$ and $\frac{1}{\sqrt{2}}(^2\nu_5 - ^2\bar{\nu}_5)$ leaving unbroken just the Standard-Model symmetry, $SU(3)_C \times SU(2)_L \times U(1)_Y$.

In principle, the ν_5 and ν_4 VEV's in Eq. (3.1) should grow along two independent flat directions of the scalar potential. We are going to discuss a different and more economic scenario where just one *almost* flat direction is sufficient to generate the two IS. This is achieved for a specific form of the $^1\nu_5 - ^2\nu_4$ interactions that will fix a *small* $< ^2\nu_4 >$ component when $< ^1\nu_5 >$ grows along the flat direction and will define a hierarchy of type $y = O(x^a)$.

To discuss the orders of magnitude, we will consider the model with superpotential

$$P = \frac{\alpha}{M_c^{2n-3}}(^1\nu_5 ^1\bar{\nu}_5)^n + \frac{\beta}{M_c^{2m-3}}(^2\nu_4 ^2\bar{\nu}_4)^m + \frac{\gamma}{M_c^{2(i+j)-3}}(^1\nu_5 ^1\bar{\nu}_5)^i(^2\nu_4 ^2\bar{\nu}_4)^j \quad (3.2)$$

and effective scalar masses destabilizing the potential along $^1\nu_5$ (i.e., $m_{^1\nu_5}^2 + m_{^1\bar{\nu}_5}^2 \equiv -m_s^2 < 0$). The actual value of the exponents in Eq. (3.2) would be fixed by the discrete symmetries of the particular manifold, while the parameters α, β and γ are arbitrary complex coefficients that we assume order 1. As shown in the previous section, if the $^1\nu_5$ self-interactions (first term in P) define an *almost* flat direction [$n = 4 - 6$ in Eq. (3.2)] these fields will develop large VEV's $< ^1\nu_5 > = < ^1\bar{\nu}_5 > = xM_c$, with F contributions in the effective potential given by

$$F_{^1\nu_5} = F_{^1\bar{\nu}_5} = n\alpha M_c^2 x^{2n-1} \approx m_s M_c x, \quad (3.3a)$$

$$F_{^2\nu_4} = F_{^2\bar{\nu}_4} = 0. \quad (3.3b)$$

Now, we study the possibility of VEV's along the nonflat [$m = 2 - 3$ in (3.2)] $^2\nu_4$ direction. If $< ^2\nu_4 > = < ^2\bar{\nu}_4 > = yM_c$, the F contributions become

$$F_{^1\nu_5} = F_{^1\bar{\nu}_5} = n\alpha M_c^2 x^{2n-1} + i\gamma M_c^2 x^{2i-1} y^{2j}, \quad (3.4a)$$

$$F_{^2\nu_4} = F_{^2\bar{\nu}_4} = m\beta M_c^2 y^{2m-1} + j\gamma M_c^2 x^{2i} y^{2j-1}. \quad (3.4b)$$

The dimensionally less suppressed contribution is the first term in $F_{^2\nu_4}$. We find, however, that if the discrete symmetries of the model fix the exponents in (3.2) such that

$$i + \frac{n}{m}j = n, \quad (3.5)$$

there will be a value of ϕ_2 (small respect to ϕ_1) that exactly cancels the nonflat contributions in $F_{2\nu_4}$ while adding to $F_{1\nu_5}$ just *almost* flat contributions (of the same order as the one introduced by ϕ_1). This value of $\phi_2(= yM_c)$ is

$$y = \left(\frac{-j\gamma}{m\beta} \right)^{\frac{1}{2(m-j)}} x^{\frac{n}{m}}, \quad (3.6)$$

and the $F_{1\nu_5}$ term fixing the minimum becomes

$$F_{1\nu_5} = [n\alpha + i\gamma \left(\frac{-j\gamma}{m\beta} \right)^{\frac{j}{(m-j)}}] M_c^2 x^{2n-1}. \quad (3.7)$$

Depending on the specific value of α , β and γ but without need of fine tuning this combination of two terms in $F_{1\nu_5}$ may be more flat than the direction with $\phi_2 = 0$, and then may correspond to a deeper minimum. Although both minima will be of the same order [$O(|F_{1\nu_5}|^2)$], they are separated (when ϕ_2 varies from zero to the value that cancels $F_{2\nu_4}$) by the barrier $|F_{2\nu_4}|^2$ in Eq. (3.4b).

We should emphasize that the ingredients to generate such a minimum [a flat direction favoured by the effective masses and interacting with another direction as specified in Eq. (3.5)] are not rare in Calabi-Yau models, where typically one is left at M_c with many different families and large groups of discrete symmetries. In Ref. 9, for example, is singled out a model (unique within its class) that satisfies certain phenomenological requirements. We find that in this case the IS have to be generated through the discussed mechanism, with $n = 6$, $m = 2$, $i = 3$ and $j = 1$; for $m_s = 1 \text{ TeV}$ and $M_c = 10^{16.5} \text{ GeV}$ we obtain

$$\phi_1 = O(10 \sqrt{\frac{m_s}{M_c}}) M_c \approx 10^{15.2} \text{ GeV}; \quad \phi_2 = O(x^3) M_c \approx 10^{12.5} \text{ GeV}. \quad (3.8)$$

Other possibilities are tabulated in Table I.

IV. LOW ENERGY IMPLICATIONS

As pointed out by Arnowitt and Nath,⁸ the flatness arguments have implications over the low-energy matter content. In particular, they found that in the Tian-Yau model the symmetry breaking along two *almost* flat directions leaves a minimum of two (and possibly four) nonchiral neutrinos with light masses (order m_s). We are going to show that our way of generating the IS has different phenomenological implications.

The nonchiral neutrinos in the two vectorlike $\lambda + \bar{\lambda}$ multiplets will receive mass contributions from gauge and chiral interactions. The first type of interactions will combine them with the gauginos associated to the lost symmetries, as stated in the previous section. Chiral interactions will affect in a different way to the fields involved in each of

the IS. While ${}^1\nu_5$, ${}^1\bar{\nu}_5$, ${}^1\nu_4$ and ${}^1\bar{\nu}_4$ will receive from terms $({}^1\lambda^1\bar{\lambda})^n$ in P order 1 TeV masses, the neutrinos in the second family will receive from $({}^2\lambda^2\bar{\lambda})^m$ and $({}^1\lambda^1\bar{\lambda})^i({}^2\lambda^2\bar{\lambda})^j$ heavy mass contributions [in the case corresponding to (3.8), the masses come from terms $({}^2\lambda^2\bar{\lambda})^2$ and are order $y^2 M_c \approx 10^9 GeV$]. Considering all these mass contributions we obtain

- a. one neutrino with mass $m_1 = O(1 TeV)$: $\frac{1}{\sqrt{2}}({}^1\nu_5 + {}^1\bar{\nu}_5) \equiv \nu_1^c$,
- b. three neutrinos with intermediate mass m_2 [$m_2 = O(10^9 GeV)$ for the case in (3.8)]: $\frac{1}{\sqrt{2}}({}^2\nu_4 + {}^2\bar{\nu}_4) \equiv \nu_2^c$, $[{}^2\nu_5 + O(y/x){}^1\bar{\nu}_4] \equiv \nu_3^c$ and $[{}^2\bar{\nu}_5 + O(y/x){}^1\nu_4] \equiv \nu_4^c$.

The neutrinos whose scalar superpartner develop a VEV can not get further masses without spoiling the flatness arguments. For this reason ν_1^c is necessarily light; if the flatness were exact, it would be massless. Although the presence of ν_1^c below 1 TeV is a model independent feature, we find that definite predictions over its observable effects would require a specific model. This neutrino does not couple in P with the rest of low-energy fields (see Section II). However, relevant couplings may appear in the gauge sector. After the first stage of breaking the gauginos $(\lambda^- \lambda^0)_L$ mix with the three chiral lepton doublets $(e \nu)_i$ ($i = 1, \dots, 3$), through mass terms of type

$$\langle {}^1\bar{\nu}_5^\dagger \rangle \begin{pmatrix} {}^1\bar{e} \\ {}^1\bar{\nu} \end{pmatrix} \begin{pmatrix} \lambda^- \\ \lambda^0 \end{pmatrix}_L ; \quad \langle ({}^1\nu_5 {}^1\bar{\nu}_5)^a \rangle \begin{pmatrix} {}^1\bar{e} \\ {}^1\bar{\nu} \end{pmatrix} \begin{pmatrix} e \\ \nu \end{pmatrix}_i . \quad (4.1)$$

Then, the gauge interactions of these gauginos with ν_1^c generate Yukawa couplings of type $({}^1\bar{e} {}^1\bar{\nu})^\dagger \bar{\nu}_1^c (e \nu)_i$. The sleptons $({}^1\bar{e} {}^1\bar{\nu})$ will have an order $\phi_2/\phi_1 \equiv \sin \theta$ component along $({}^2\tilde{h}' - {}^2\tilde{h}'^0)$, which may have sizeable components along the space of the three low-energy sleptons. It results

$$L_Y = \lambda_{ij} \sin \theta \begin{pmatrix} \tilde{e} \\ \tilde{\nu} \end{pmatrix}_i^\dagger \bar{\nu}_1^c \begin{pmatrix} e \\ \nu \end{pmatrix}_j + h.c. \quad (4.2)$$

All the mixings are allowed by the Z_2 and Z_3 matter parities¹⁰ defined from discrete symmetries of the Tian-Yau manifold (see Table II). The Yukawas in Eq. (4.2) would induce at one-loop level effective operators of type anomalous magnetic moment (l_i , with $i = 1, 2$ and 3 , stands for the charged leptons e , μ and τ)

$$L_{eff} = \frac{e}{4m_j} F^{\mu\nu} \bar{l}_i \sigma_{\mu\nu} (a_L^{ij} P_L + a_R^{ij} P_R) l_j + h.c. \quad (4.3)$$

producing lepton-number violation. For instance, in the process $\mu \rightarrow \gamma e$ one obtains an amplitude proportional to $m_\mu/\bar{m}^2$, where $\bar{m} = O(m_{slepton} + m_1)$, with $a_L = (m_e/m_\mu)a_R$ and then an emitted electron almost 100% left-handed polarized. Although this is an interesting effect, very constrained experimentally [$|a_L^{12}| + |a_R^{12}| \leq O(10^{-13})$]^{16,17}, its size here depends on ratios of nonrenormalizable terms and appears to be, in consequence, model dependent.

The fields in the intermediate region m_2 could be relevant for the masses of the three chiral neutrinos ν_i . Since the scalars ${}^2\nu_5$ and ${}^2\bar{\nu}_5$ do not develop VEV's, there could be sizeable couplings in P of type $h^0 \nu_i \nu_{3,4}^c$ inducing see-saw neutrino masses. (Note that trilinears of $\nu_{3,4}^c$ with the two Higgs doublets are forbidden by the matter parities in Table II.) The heavy fields $\nu_{3,4}^c$ could make massive through this mechanism two of the three chiral neutrinos. Besides, one expects the couplings of ν_4^c (mainly in a $\bar{27}$ of E_6) to be smaller than those of ν_3^c , what suggests one chiral neutrino (ν_h) heavier than the other two. In this scenario, cosmological constraints¹⁸ imply $m_{\nu_h} < O(100 \text{ eV})$. Assuming typical leptonic couplings, [$O(m_\tau / \langle h^0 \rangle)$] this translates into a condition for the heavy neutrinos, $m_2 > O(m_\tau^2 / 100 \text{ eV}) = O(10^7 \text{ GeV})$, which is fulfilled just in some of the models in Table I. It is possible, however, to make $\nu_{3,4}^c$ heavier and/or to suppress their couplings with ν_i . Further masses for $\nu_{3,4}^c$ would arise through their combination with other neutrinos in 27 and $\bar{27}$ multiplets and with the E_6 singlets present in Calabi-Yau models. These mass terms would not spoil the flatness in the scalar potential (as it would happen with masses for ν_1^c and ν_2^c). The second possibility (suppression of the couplings) is also quite model dependent. Actually, the Z_3 matter parity forbids this simplified scheme of see-saw masses. In this case, after diagonalizing the 5×5 mass matrix corresponding to ν_i , ν_3^c and ν_4^c , we find two heavy neutrinos [ν_4^c and $\nu_3^c + O(m_\tau/m_2)\nu_i$] defining a Dirac field plus three massless neutrinos [$\nu_i + O(m_\tau/m_2)\nu_3^c$]. For these reasons, although the scenario suggests masses for the Standard Model neutrinos, even the estimation of orders of magnitude precises a more specific model.

V. CONCLUSIONS

We have analyzed a mechanism that requires just one *almost* flat direction in the scalar potential to generate the two stages of gauge-symmetry breaking precised in rank-six Calabi-Yau models. This mechanism is not contained in the general analysis done by Arnoult and Nath (where two flat directions are assumed), and consequently its implications are different from those. We find the second scale of breaking several orders of magnitude smaller than the first. We also find only one nonchiral neutrino (the one along the flat direction) necessarily light, with mass contributions in the intermediate region for the neutrinos in the multiplet defining the second scale. This scenario has non-Standard low-energy implications, such as see-saw masses for the chiral neutrinos and lepton-number violating processes. However, their size appears to be quite model dependent. An analysis of an explicit model where this mechanism is realised will be done elsewhere.

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Table I. Order of magnitude of the IS and the masses of the neutrinos in the multiplet defining ϕ_2 .

n	ϕ_1/M_c	m	ϕ_2/M_c	(i, j)	$m_2(\text{GeV})$
6	$10^{-1.3}$	2	$10^{-4.0}$	(3, 1)	10^9
		3	$10^{-2.7}$	(2, 2)	10^7
		3	$10^{-2.7}$	(4, 1)	10^7
		4	$10^{-2.0}$	(3, 2)	10^5
4	$10^{-2.7}$	2	$10^{-5.4}$	(2, 1)	10^6

Table II. Z_2 and Z_3 matter parities defined from the discrete symmetries of the Tian-Yau manifold.¹⁰

Field	Z_2	Z_3	Field	Z_2	Z_3
$(e \nu)_i$	-1	1	$< {}^1\nu_5 >, < {}^1\bar{\nu}_5 >$	+1	1
e^c_i	-1	α^2	$< {}^2\nu_4 >, < {}^2\bar{\nu}_4 >$	+1	1
$(u d)_i$	-1	1	${}^1\nu_4$	-1	α
u^c_i	-1	α	${}^1\bar{\nu}_4$	-1	α^2
d^c_i	-1	α^2	${}^2\nu_5$	-1	α^2
$(h^0 h^+)$	+1	α^2	${}^2\bar{\nu}_5$	-1	α
$(h'^- h'^0)$	+1	α	$({}^1e {}^1\nu), \overline{({}^1e {}^1\nu)}$	-1	1
ν_1^c	+1	1	$({}^2h'^- {}^2h'^0), \overline{({}^2h'^- {}^2h'^0)}$	-1	1

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