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CAN THE SUPERSTRING INSPIRE THE STANDARD MODEL ?

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A B S T R A C T

We discuss general features of models in which the $E_8 \times E_8$ heterotic superstring is compactified on a specific Calabi-Yau manifold. The gauge group of rank-6 in four dimensions is supposed to be broken down at an intermediate scale m_I to the standard model group $SU(3)_C \times SU(2)_L \times U(1)_Y$, as a result of two neutral scalar fields acquiring large vacuum expectations (vev's) in one of many flat directions of the effective potential. We find that it is difficult to generate such an intermediate scale by radiative symmetry breaking, whilst such models have prima facie problems with baryon decay mediated by massive particles and with non-perturbative behaviour of the gauge couplings, unless $m_I \gtrsim 10^{16}$ GeV. Rapid baryon decay mediated by light particles, large neutrino masses, other $\Delta L \neq 0$ processes and flavour-changing neutral currents are generic features of these models. We illustrate these observations with explicit calculations in a number of different models given by vev's in different flat directions.

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1. - INTRODUCTION

As the superstring¹⁾ is rapidly winning over the hearts and minds of high energy theorists, it may soon live up to the claim that the superstring is indeed the greatest invention since sliced bread. To be sure, there remains a great challenge to make the connection between the theory of everything (TOE) provided by the heterotic superstring²⁾ at energies $0(10^{19})$ GeV in ten dimensions, and lower-energy physics in four dimensions. The latter seems to be well described by the standard model $SU(3)_C \times SU(2)_L \times U(1)_Y$, so the following questions arise. Can the standard model be obtained at high energies directly from the TOE in a natural way? or should one expect that physics below $0(1)$ TeV contains additional physics beyond the standard model?

There is already a considerable amount of literature on superstring phenomenology or superstring-inspired models³⁾, but as yet very little on true superstring models⁴⁾. The distinction between inspired and true models is that in the former, one merely hopes that a compactification manifold exists with the right properties, such as its Euler characteristic and discrete symmetries, so that the effective low-energy theory can produce acceptable phenomenology. In a true model, one does not merely hope or ask for a specific manifold and discrete symmetry. One starts with a given manifold, and symmetries are either there or they are not. Only one example of a true model has been worked out in any detail⁴⁾, which is based on the Yau⁵⁾ manifold $CP^3 \times CP^3$ with some important extra assumptions. In this paper we examine some general features of models based on this Yau manifold, such as baryon- and lepton-number-violating processes, fermion masses and flavour-changing neutral interactions.

Even within the class of true models, it is not a straight path from the TOE to the standard model. The first fork in the road is how to compactify the six surplus dimensions. Compactifications on Calabi-Yau manifolds yield four-dimensional gauge groups which are rank-5 or 6 subgroups of E_6 ⁶⁾. Alternative manifolds have been proposed which might yield a rank 4 gauge group⁷⁾, but these seem to suffer from other theoretical and phenomenological problems⁸⁾. Brighter prospects may be offered by intrinsically stringy compactifications on orbifolds⁹⁾, or by string theories formulated directly in four dimensions¹⁰⁾, but no appealing phenomenological example has yet been developed. Moreover, many such stringy compactifications tend to give larger four-dimensional gauge groups¹¹⁾ rather than smaller ones.

Whichever option we took at the previous fork in the road, we are faced with another fork: how to break the gauge group down from higher rank to the rank-4 standard model. The tools available to break the rank-5 or 6 gauge group left over from Calabi-Yau compactification are two $SU(3)_C \times SU(2)_L \times U(1)_Y$ singlet fields which we label v^c and N . The former is an $SU(5)$ singlet in a 16 of $SO(10)$, whilst the latter is an $SO(10)$ singlet in the fundamental 27 of E_6 . To break down from rank 5 (6), one (both) of these fields must acquire a vacuum expectation value. The $CP^3 \times CP^3$ manifold⁵⁾ considered here yields initially the rank-6 gauge group $SU(3)_C \times SU(3)_L \times SU(3)_R$, and hence both v^c and N must have non-vanishing vev's. If there is to be no new physics beyond the standard model at scales less than $O(1)$ TeV then both fields must have vev's at some intermediate scale $\langle v^c, N \rangle \sim M_I \gg O(1)$ TeV.

Such a model was first proposed in Ref. 4). The intermediate scales were generated along a certain flat direction of the potential followed by a specific chain of symmetry breaking, and finally a low energy spectrum was extracted. Here, we will make a broad survey of such models using various different flat directions. We find that these models in general have problems with flavour-changing neutral interactions¹²⁾, fermion masses and baryon- and lepton-number-violating interactions, which can only be avoided by a quite specific, and otherwise unmotivated, choice of flat direction.

In Section 2, before discussing these issues, we first review the (im)possibility of generating an intermediate scale using radiative symmetry breaking¹³⁾. In addition to this difficulty, we find problems with non-perturbative behaviour of the gauge couplings and with baryon decay unless $M_I > O(10^{16})$ GeV, while at the same time we find problems with an excess entropy production at the intermediate-scale phase transition unless $M_I < O(10^{12})$ GeV^{13),14)}. We also preview why we believe that models with both $\langle v^c \rangle \neq 0$ and $\langle N \rangle \neq 0$ should a priori violate maximally baryon and lepton number. Other problems with the model of Ref. 4) are discussed in Ref. 15).

In Section 3, we survey several flat directions in the $CP^3 \times CP^3$ model. For the sake of argument, we ignore (blindly) the problems of the previous section and calculate the allowed Yukawa couplings and the light particle spectrum.

In Section 4, we describe the general features of these models such as the possibility of generating fermion masses without excessive flavour-changing neutral interactions, the generation of Higgs doublet mixing, which is

essential in the standard model, and $\Delta B \neq 0$ and $\Delta L \neq 0$ interactions. Our summary and conclusions are given in Section 5.

2. - INTERMEDIATE SCALE MODELS?

If post-compactification symmetry breaking is to occur in such a way as to yield the standard model at a scale above $O(1)$ TeV, it is necessary to establish an intermediate scale in the model. Besides giving the standard model at energies $\gg 1$ TeV, intermediate scales have also been proposed as a means to generate small but finite neutrino masses using non-renormalizable interactions¹⁶⁾, generate a baryon asymmetry¹⁷⁾, and serve other useful phenomenological purposes.

Although vacuum expectation values of the singlet fields v^c or N can in principle be quite large, their affinity with conventional squark, slepton and Higgs fields implies that the supersymmetry-breaking parameters associated with these fields must be kept small $\sim m_W$. Because of the form of the scalar potential the global minimum would occur when all fields lay at $\langle \phi \rangle = 0$ were it not for supersymmetry breaking. If supersymmetry breaking is due to gaugino condensation in the hidden sector¹⁸⁾, gravitational-strength couplings feed this breaking into the observable sector resulting in softly broken supersymmetric masses for scalars and gauginos, and gauge symmetry breaking becomes possible¹⁹⁾.

The general form of the scalar potential can be written as

$$V(\phi) = m_0^2 \phi^2 + (ag^2 + b\lambda^2) \phi^4 + \left(\begin{smallmatrix} \text{higher} \\ \text{order} \\ \text{terms} \end{smallmatrix} \right) \quad (1)$$

where ϕ is a symmetry-breaking order parameter, which takes values along a ray in the group space of v^c and N , $|m_0^2| \lesssim O(m_W^2)$ and ag^2 and $b\lambda^2$ represent contributions from the D and F terms^{*}). The squared mass m_0^2 is determined by renormalization group equations such as

^{*}) There could in addition be supersymmetry-breaking quartic terms with coefficients $O(m^2/m_P^2)$, but these do not affect the determination of vev's $\langle 0|\phi|0 \rangle \sim O(m_P)$, and are ignored hereafter.

$$8\pi^2 \frac{dm_0^2}{dt} = -C \cdot g^2 m_{1/2}^2 + C' \lambda^2 m_0^2 \quad (2)$$

where $t \equiv \ln(Q/m_p)$, C and C' are positive, model-dependent numbers $O(1)$, $m_{1/2}$ is a generic gaugino mass, m_0 a generic scalar mass, g a gauge coupling and λ a Yukawa coupling. If at the compactification scale $m_0^2 \sim 0$ ¹⁹⁾, the gauge couplings will initially drive $m_0^2 > 0$, and then the Yukawa couplings (if present) could drive $m_0^2 < 0$ and allow non-zero vacuum expectation values for ϕ .

If indeed the renormalization group equations (2) lead to $m_0^2 < 0$ at some low renormalization scale Q_0 , the potential (1) would typically give $\langle \phi \rangle \sim \sqrt{-m_0^2} \sim m_W$ rather than an intermediate scale. However, the D-term $ag^2\phi^4$ in (1) originates from a more general expression $(\phi_i^\dagger T_\alpha \phi_i - \bar{\phi}_i^\dagger T_\alpha \bar{\phi}_i)^2$, and it is possible for this term to vanish if a second field $\bar{\phi}$ also picks up a vacuum expectation value $|\langle \phi \rangle| = |\langle \bar{\phi} \rangle|$ ²⁰⁾. To make all the ϕ^4 terms vanish, one must require F-flatness so that $b\lambda^2\phi^4$ vanishes. A quick inspection of the E_6 invariant $(\underline{27})^3$ superpotential couplings shows that there are no cubic superpotential terms $\alpha(v^c, N)^2$ and hence the renormalizable terms in the potential can be both F flat as well as D flat.

In order to determine $\langle \phi \rangle$, one must now look at the non-renormalizable couplings in the superpotential. To fourth order, $f = (\underline{27})^2(\overline{27})^2$ and the scalar potential is of the form

$$V = m_0^2 \phi^2 + \lambda' \frac{\phi^2 \bar{\phi}^2}{M_c^2} (\phi^2 + \bar{\phi}^2) + \left(\begin{array}{c} \text{higher} \\ \text{order} \\ \text{terms} \end{array} \right) \quad (3)$$

where M_c is the compactification scale. Along the D-flat direction ($\phi = \bar{\phi}$) this yields $\langle \phi \rangle^2 \sim M_c \sqrt{-m_0^2} \sim (0(10^{10}) \text{ GeV})^2$. If the potential is also flat at fourth order in ϕ , it would be necessary to consider still higher-order terms. For instance, if the first non-flat contribution is sixth order in f , this will yield $\langle \phi \rangle^4 \sim M_c^3 \sqrt{-m_0^2}$.

Other possible contributions to the potential are terms involving singlets such as $S(\underline{27})(\overline{27})$ ⁷⁾. There are many possible singlet fields S not all of which couple to $(\underline{27})(\overline{27})$. For example, there are singlet fields associated with deformations of the compactified manifold, K , whose properties can be calculated from the cohomology group $H^{(2,1)}(K)$. These singlets do not couple to $(\underline{27})(\overline{27})$ ⁶⁾. Singlets for which there could in principle be a coupling present are the fields

transforming as $(\underline{1}, \underline{8})$ under $E_6 \times SU(3)$ ⁷⁾ but there is no general method for calculating the number of these singlets nor their couplings to the $(\underline{27})(\overline{\underline{27}})$. Moreover, arguments have been given ²¹⁾ that these $(\underline{1}, \underline{8})$ fields acquire masses $O(m_p)$ from non-perturbative effects, in which case they are absent from the low-energy effective theory. At the present time, we feel it is preferable to neglect the possible rôle of singlet fields.

The general strategy for producing an intermediate scale is then as follows: 1) Look for a flat direction; 2) use the renormalization group equations to drive m_0^2 negative at a renormalization scale Q_0 large enough to accommodate a large intermediate scale vev $\langle \phi \rangle \lesssim Q_0$. To see if this second step is actually feasible, we consider the renormalization group equation (2). Since we will be studying behaviour in a restricted range of t , we approximate $g \approx \text{constant} \equiv \bar{g}$ and also take $m_{\frac{1}{2}} \propto g^2$ as a constant. We expect m_0^2 to be driven negative most quickly by a large Yukawa coupling λ , typically after a few orders of magnitude in $\mu = e^t$. These reach very rapidly an infra-red fixed point given by

$$\bar{\lambda}^2 \simeq a \bar{g}^2 \simeq \text{constant} \quad (4)$$

where a is a number of $O(1)$ depending on the particle content. The fastest possible rate of growth of a squared scalar mass m_0^2 is

$$8\pi^2 \frac{dm_0^2}{dt} = -A_0^2 \bar{g}^2 m_{\frac{1}{2}}^2 \Rightarrow m_0^2 = \frac{A_0 \bar{g}^2 m_{\frac{1}{2}}^2}{8\pi^2} |t| \quad (5)$$

where $A_0 = O(1)$. Thus the fastest possible rate by which some other m_0^2 can be driven negative is

$$8\pi^2 \frac{dm_0^2}{dt} = -C \bar{g}^2 m_{\frac{1}{2}}^2 + A_0 \frac{C' a \bar{g}^4 m_{\frac{1}{2}}^2}{8\pi^2} |t| \quad (6)$$

This yields

$$m_0^2 = \left(C - \frac{a A_0 C' \bar{g}^2 |t'|}{16\pi^4} \right) \frac{|t'| \bar{g}^2 m_{\frac{1}{2}}^2}{8\pi^2} \quad (7)$$

where $t' \equiv \ln Q/m_s$ with m_s the mass scale at which supersymmetry breaking first appears: $m_s \sim$ the gaugino condensation scale Λ_c in one favoured scenario ¹⁹⁾. The expression (7) vanishes at

$$|t'_0| = \left(\frac{c}{aA_0c'}\right)\left(\frac{16\pi^2}{\bar{g}^2}\right) \quad (8)$$

which gives a lower bound on the hierarchy m_s/Q_0 . We take as typical values

$$c = \frac{5}{3}, \quad A_0 = 5, \quad c' \approx 3, \quad a = \frac{8}{3}, \quad \text{and} \quad \bar{g}^2 = 4\pi\alpha_{4VT} = 1 \quad (9)$$

for ν^c or N

in which case $|t'_0| \approx 7$ corresponding to $Q_0 = 10^{-3} m_s$: in the context of our assumptions there is no way of driving m_0^2 negative above this scale. Even if $m_s \approx 10^{18}$ GeV, which is excluded in most models based on gaugino condensation, we have the upper bound

$$m_I = \langle \phi \rangle \ll Q_0 \lesssim 10^{-3} m_s \lesssim 10^{15} \text{ GeV} \quad (10)$$

This bound is very conservative, and realistic cases generally have $m_I \lesssim 10^{-6} m_s$ and $m_s \sim 10^{-2} m_P$ so that $m_I \lesssim 10^{10}$ GeV. We will see shortly that values of m_I larger than these are required by other phenomenological considerations.

Despite the problem that large intermediate scales cannot easily be generated by radiative symmetry breaking, one could simply hope that some other (unknown) mechanism is responsible for m_I in which case one might abandon relations such as $m_I^{n-2} \sim m_W M_C^{n-3}$ and let m_I take any value between m_W and M_C . There are, however, other constraints on m_I which we now discuss. Proton decay can be mediated by extra charge $-1/3$ colour triplet particles^{12),20)}, called here D quarks, and by D^c antiquarks. Because their masses will typically be $O(\lambda m_I) \leq O(m_I)$, constraints due to proton stability on dimension-5 and 6 operators require $m_I > O(10^{16})$ GeV²²⁾. In addition, if we consider the renormalization group evolution of the strong gauge coupling α_s above m_I , non-perturbative gauge effects can take over unless $m_I > O(10^{16})$ GeV as well¹³⁾.

At the same time, it is possible to argue that unless $m_I < O(10^{12})$ GeV, the phase transition producing $\langle \phi \rangle \sim m_I$ will generate an excess of entropy thus washing away the baryon asymmetry of the Universe^{13),14)}. If one considers the high temperature corrections to $V(\phi)$, for $\phi \ll T < m_I$, $V_T(\phi) \sim \phi^2 T^2$ so that the phase transition does not occur until $T^2 \sim -m_0^2$ (tunnelling is found to be negligible). Because the potential is very flat and ϕ is coupled directly only to fields with masses $\sim m_I$, ϕ decays very slowly and the cosmological evolution of $\langle \phi \rangle$ becomes a damped oscillation about $\langle \phi \rangle \approx m_I$. The entropy release Δ was

calculated¹⁴⁾ to be $\Delta \approx 0(10^2)m_I^3/m_0^{5/2}m_P^{1/2}$. For $\Delta < 10^6$, one has the limit $m_I < 0(10^7) \text{ GeV}^{14)}$. This limit is relaxed somewhat¹³⁾ if one produces the baryon asymmetry at low temperatures via the Affleck-Dine mechanism²³⁾. In this case we found that in order to maintain $n_B/S \gtrsim 10^{-11}$ one must have $m_I \lesssim 0(10^{12}) \text{ GeV}^{13)}$. This is in apparent conflict with the previous bounds on m_I .

Even if one could avoid all the above difficulties, we expect several other problems to appear²⁴⁾ when both $\langle v^c \rangle$ and $\langle N \rangle$ are non-zero, as is necessary in order to obtain the standard model at an intermediate scale from a rank-6 model. If only one of $\langle v^c \rangle$, $\langle N \rangle \neq 0$, taken by convention to be $\langle N \rangle$, then one may hope that all the D and D^c states will be heavy, while all the light charge $[1/3]$ states are d or d^c states. In this case, as is well known^{12),19),24)}, there are no $\Delta B \neq 0$ interactions in the effective low-energy theory and there is no L violation due to $\langle v^c \rangle$. However, if both $\langle v^c \rangle$ and $\langle N \rangle \neq 0$, the light charge $+1/3 \bar{3}$ states are in general mixtures of d^c and D^c , and $\Delta B \neq 0$ interactions are possible. Moreover, $\langle v^c \rangle \neq 0$ means that the vacuum has $\Delta L \neq 0$. In general, these $\Delta B \neq 0$ and $\Delta L \neq 0$ effects would be maximal, as we will discuss in Section 4 after first setting up the general structure of low-energy models in Section 3.

3. - LOW ENERGY MODELS

Despite the problems raised in the previous section for models utilizing intermediate scales, we will in this section look at a sample of models, based on the manifold $CP^3 \times CP^3$ ⁵⁾ as was analyzed in Ref. 4). Although we do not know of a convincing mechanism for generating an intermediate scale, we believe nevertheless that intermediate scales require flat directions in the scalar potential, as described in the previous section. If we consider generically a typical scalar potential as in Eq. (1) with $\lambda \sim 1$ and $m_0^2 \sim -m_W^2$ we see that $m_I \sim [m_W m_c^{n-3}]^{1/n-2}$. Thus to get $m_I \gtrsim 10^{16} \text{ GeV}$ would seem naively to require $n \gtrsim 9$, i.e., the superpotential is flat up to ninth order! As will become apparent shortly, this degree of flatness is not likely to be feasible in the types of models that we will discuss. As in Ref. 4), we consider models which are flat at least to sixth order ($n=6$), corresponding to an intermediate scale $m_I \gtrsim 0(10^{14}) \text{ GeV}$.

We begin by discussing the relevant properties of the $CP^3 \times CP^3$ model⁵⁾ and for the most part follow the notation of Ref. 4). The model is based on the

manifold defined as the subspace R_0 of $CP^3 \times CP^3$ on which polynomials in the homogeneous co-ordinates (x,y) of the two CP^3 's with bidegrees $(3,0)$, $(0,3)$ and $(1,1)$ vanish. The discrete symmetry group acting on R_0 , necessary for flux breaking, is chosen to be Z_3 , which is taken to be embedded in E_6 in such a way as to leave unbroken the gauge group $[SU(3)]^3$. Decomposing a $\underline{27}$ of E_6 under $SU(3)_C \times SU(3)_L \times SU(3)_R$: $\underline{27} = (\underline{1}, \underline{3}, \underline{3}) + (\underline{3}, \underline{3}, \underline{1}) + (\underline{3}, \underline{1}, \underline{3})$, it was found that the massless modes in R_0/Z_3 consist of nine lepton generations λ_i and seven quark and antiquark generations q_i and Q_i along with six mirror lepton generations $\bar{\lambda}_i$ and four mirror quark and antiquark generations $\bar{q}_i$ and $\bar{Q}_i$. The symmetry properties of these fields are tabulated in Ref. 4), along with the allowed dimension-four superpotential terms. In Table 1, we list the allowed cubic terms of the superpotential and include from Ref. 4) the allowed quartic terms P_i .

3.1 - F-flatness

In order to produce the standard model at a large intermediate scale, it is necessary to reduce the rank by two units using vev's for both the v^c and the N fields. We will use the notation that under an $SO(10)$ decomposition, the $\underline{27}$ of E_6 consists of a $\underline{16}$ containing the ordinary $SU(5)$ $\underline{10}$ and $\bar{\underline{5}}$ plus an $SU(5)$ singlet v^c ; a $\underline{10}$ containing $SU(5)$ multiplets $\underline{5}$ (D, H) and $\bar{\underline{5}}$ (D^c, H^c), where (D, D^c) are colour triplets and (H, H^c) are $SU(2)_L$ doublets, v^c and an $SO(10)$ singlet N . In order to have a chance to produce large vev's for N and v^c , we check for flat directions through sixth order in the scalar potential, which means at least through fourth order in the superpotential.

A generic fourth order superpotential term for lepton superfields can be written as $\lambda_a \lambda_b \bar{\lambda}_c \bar{\lambda}_d$, and F-flatness requires that

$$\left| \frac{\partial f}{\partial \lambda_a} \right|^2 = \left| \frac{\partial f}{\partial \lambda_b} \right|^2 = \left| \frac{\partial f}{\partial \bar{\lambda}_c} \right|^2 = \left| \frac{\partial f}{\partial \bar{\lambda}_d} \right|^2 = 0 \quad (11)$$

Let us consider a few examples of specific directions to see whether such flatness is possible. We will assume as in Ref. 4) that the first stage of breaking occurs by giving vev's to v_1^c , $\bar{v}_5^c$ and $\bar{v}_6^c$ to break $[SU(3)]^3$ down to $SU(3)_L \times SU(2)_L \times SU(2) \times U(1)$, although many other possibilities exist. Breakdown to the standard model will then require one or more N and $\bar{N}$ field to acquire a vev.

If for example we consider directions with $\langle N_1 \rangle$, $\langle \bar{N}_1 \rangle$ and $\langle \bar{N}_2 \rangle$ non-zero, we find that the only possible contributions to the potential come from terms in P_2 , namely $N_1 N_3 \bar{N}_{1,2}^2$ where $\bar{N}_{1,2}^2$ represents a quadratic $a\bar{N}_1^2 + b\bar{N}_1 \bar{N}_2 + c\bar{N}_2^2$ in $\bar{N}_1$ and $\bar{N}_2$. Differentiating with respect to N_3 yields a non-zero contribution to the scalar potential V , unless the ratio of $\bar{N}_1$ to $\bar{N}_2$ is such that this quadratic also vanishes (recall that D-flatness requires $\langle N_1^2 - \bar{N}_1^2 - \bar{N}_2^2 \rangle = 0$). As this is the only potential problem for F-flatness there remains an F-flat direction for some combination of N_1 and $(\bar{N}_1, \bar{N}_2)$ denoted by $1(\bar{1}, \bar{2})$.

A second example with more problems is the $N_8, \bar{N}_1, \bar{N}_2$ [denoted by $8(\bar{1}, \bar{2})$] direction. In this case, potential problems arise from the following terms: (1) $P_3: v_1^c N_8 \bar{v}_{5,6}^c \bar{N}_{5,6}$ where differentiating with respect to $\bar{N}_{5,6}$ and the vev's for $v_1^c, v_{5,6}^c$ leave $V \propto |N_8|^2$ creating two problems in all $8(\bar{1}, \bar{j})$ directions; and (2) $P_9: N_8 N_9 \bar{N}_{1,2}^2$. Clearly, the P_3 mass terms cannot be made to vanish simultaneously in general, although the P_9 mass terms can.

In Table 2, we list the quartic superpotential terms which contribute to the non-flatness of V in the directions $N_i \bar{N}_j$ assuming $\langle v_1^c \rangle \langle v_5^c \rangle \langle v_6^c \rangle$ are all non-zero. We can only be sure of F-flatness when one or less entries are present, thus leaving $1(\bar{1}, \bar{2})$, $2(\bar{1}, \bar{2})$, $4(\bar{1}, \bar{2})$, $5(\bar{1}, \bar{2})$, $6(\bar{1}, \bar{2})$, $7(\bar{1}, \bar{2})$ and $9(\bar{1}, \bar{2})$ as possible flat directions. We now exhibit the light ($m \ll m_I$) fermion spectrum corresponding to each of these flat directions.

3.2 - Light fermion spectrum

1) Quark mass matrices

Because $CP^3 \times CP^3$ models contain 11 charge $\pm 2/3$ quarks and 22 charge $\pm 1/3$ quarks, there is the possibility of a great deal of mixing. To be sure, not all of the 22 charge $\pm 1/3$ "down" quarks mix. Fields with similar standard model quantum numbers are the d^c , D^c and $\bar{D}$ antiquarks, and the D , $\bar{d}^c$ and $\bar{D}^c$ quarks. Thus, for each choice of vev's we can write down an 18×15 down quark mass matrix, which yields at least three massless charge $-1/3$ mixtures, plus possibly some additional states accompanied by mirror particles. In Table 3 we give as an example the mass matrix in the $1(\bar{1}, \bar{2})$ model. In this matrix, entries labelled C are masses coming from cubic terms in the superpotential, whilst entries Q originate from quartic terms. The cubic entries arise from the following terms in the superpotential:

$$\left. \begin{aligned} & \lambda_1 (q_2 Q_7 + q_3 Q_4 + q_4 Q_3 + q_6 Q_6 + q_7 Q_2) \\ & \bar{\lambda}_{1,2} (\bar{q}_1 \bar{Q}_2 + \bar{q}_2 \bar{Q}_1 + \bar{q}_4 \bar{Q}_4 + \bar{q}_2 \bar{Q}_3 + \bar{q}_3 \bar{Q}_2 + \bar{q}_4 \bar{Q}_1 + \bar{q}_3 \bar{Q}_4 + \bar{q}_4 \bar{Q}_3) \end{aligned} \right\} \quad (12)$$

Vacuum expectation values for v^c fields yield Dd^c mass terms, whilst those for N fields yield DD^c mass terms, and analogously for the mirror fields. The quartic terms contributing to the mass matrix are of the form $\lambda_a \bar{\lambda}_b (q_c \bar{q}_d \text{ or } Q_c \bar{Q}_d)$. For the pattern of vev's studied here, the following quartic superpotential terms give mass terms:

$$\left. \begin{aligned} & N_1 \bar{N}_{1,2} q_1 (\bar{q}_1 + \bar{q}_3) \quad , \quad N_1 \bar{N}_{1,2} q_5 (\bar{q}_2 + \bar{q}_4) \\ & \nu_1^c \bar{\nu}_{5,6}^c q_2 (\bar{q}_1 + \bar{q}_3) \quad , \quad \nu_1^c \bar{\nu}_{5,6}^c q_6 (\bar{q}_2 + \bar{q}_4) \end{aligned} \right\} \quad (13)$$

as well as analogous terms with $q_i \rightarrow Q_i$ and $\bar{q}_j \rightarrow \bar{Q}_j$.

The full 18×15 down quark mass matrix can in general be broken down to several smaller independent submatrices. Each submatrix contributes a number of light states equal to the difference between the numbers of its horizontal and vertical entries. In the $1(\bar{1}, \bar{2})$ example shown in Table 3 the full matrix breaks down to one 14×13 submatrix and two 2×1 submatrices. The resulting light antidown quarks are therefore found to be one out of each of the following combinations of fields:

$$1/(d_3^c, D_3) \quad , \quad 1/(d_4^c, D_4) \quad , \quad 1/(d_{1,2,5,6,7}^c, D_{1,2,5,6,7}^c, \bar{D}_{1,2,3,4}) \quad (14)$$

It is easier to find the light u , u^c and d quarks, which are deduced from the quartic terms above to be

$$q_{3,4,7} \quad , \quad u_{3,4,7}^c \quad (15)$$

where q_i stands for the doublet (u_i, d_i) .

2) Lepton mass matrices

The lepton mass matrices are treated similarly to the charge $\pm 1/3$ quarks. Mixing occurs between the doublet fields L , H^c and $\bar{H}$ [where $L \equiv (\lambda, \nu_\lambda)$ and $\bar{H}$ is the mirror of H], and between the doublets H , $\bar{L}$ and $\bar{H}^c$. In addition, there are two $SU(2)_L$ -doublet gaugino states, one with positive and one with negative $U(1)_Y$ hypercharge, which mix with these lepton and Higgs states. Thus we have in general a 25×22 mass matrix for the $SU(2)_L$ doublet "leptons" which can be broken up like that for the d^c above, and produces at least three light doublets. In the absence of $SU(2)_L$ breaking, these doublets do not mix with the singlet charged "leptons" which include the e^c , $\bar{e}^c$, two positively and two negatively charged gauginos. These have an 11×8 mass matrix which produces at least three light singlets. Similarly to entries in the quark mass matrix, cubic terms arise from the following couplings:

$$\lambda_1 (\lambda_1^2 + \lambda_3 \lambda_6 + \lambda_3 \lambda_7 + \lambda_4 \lambda_9) \quad (16)$$

whilst the subscripts on the quartic terms in the matrix of Table 3b refer to the superpotential terms listed in Table 1b. In this case, the 25×22 matrix breaks down into one 16×16 , two 2×1 , four 1×0 and two 0×1 submatrices, leaving the following light charged states:

$$\left. \begin{array}{l} L_2, L_5, H_2^c, H_5^c, 1/(L_4, H_4^c), 1/(L_9, H_9^c) \\ H_2, H_5, 1/(\bar{L}_{5,6}) \end{array} \right\} \quad (17)$$

In addition, the quartic terms give mass to many of the remaining singlet lepton fields e^c , ν^c , N , leaving only

$$e_{2,4,5,6,7,9}^c ; 2/(\bar{e}_{1,2,3,4}^c) ; \nu_{5,6,7}^c \quad (18)$$

and a negatively charged gaugino mixture as light states.

The same approach can be extended to the other flat $n(\bar{1}, \bar{2})$ models, obtaining the other light particle spectra exhibited in Table 4. We do not present here explicit results for the direction chosen in Ref. 4). There vev's in the $\nu_1^c - \nu_2^c$, $\bar{\nu}_5^c$, $\bar{\nu}_6^c$ and $N_8 + N_9$, $\bar{N}_1$, $\bar{N}_2$ directions were chosen. Results would be qualitatively similar to the $9(\bar{1}, \bar{2})$ model shown in Table 4. However, as stated earlier, we

disagree with the way singlet fields were treated in Ref. 4). Also, the resultant light particle spectrum was only purged of phenomenological problems by assuming ad hoc additional vev's. In particular, the matter parity which guarantees baryon stability is rather arbitrarily imposed.

There are several features of the flat models in Table 4 which appear quite general. (1) There are no N fields present in the light spectra. This complicates the realization of the HH^c mixing which is necessary for weak gauge symmetry breaking. (2) There is an abundance of mixing between the charge $+1/3$ "anti-quarks" d^c and D^c which tends to make proton decay quite rapid. This mixing is directly traceable to the assumption of simultaneous vev's for both v^c and N fields. These and other phenomenological difficulties will be studied in more detail in the next Section.

4. - PROBLEMS

We now examine the possibility of obtaining a phenomenologically reasonable low-energy model from among the multitude of possible light state spectra in the various flat directions. Of primary importance are the generation of fermion masses through the standard Higgs mechanism of electroweak gauge symmetry breaking, and HH^c mixing. We also examine the possible presence of flavour-changing neutral currents. Finally, we look for possible baryon- and lepton- number-violating interactions. For each of these potential problems, we discuss in some detail the situation in the $1(\bar{1}, \bar{2})$ model, and discuss the corresponding situations in the other models in a more sketchy way. The conclusions are all displayed in Table 5.

4.1 - Fermion masses

In general, fermion masses may arise from the following sets of superpotential terms: m_u from Hqu^c ; m_d from $H^c qd^c$ and /or LqD^c , depending whether $\langle H^c \rangle \neq 0$ and/or $\langle L \rangle \neq 0$; m_e from $H^c Le^c$; m_ν from HLv^c .

In the $1(\bar{1}, \bar{2})$ model there are two possibilities for a non-zero vev for a Higgs field H : $\langle H_2 \rangle \neq 0$ and/or $\langle H_5 \rangle \neq 0$. We see from Table 1a that H_5 has no couplings to any of the light u quarks, and hence does not contribute to any m_u ,

whereas H_2 couples to $u_2 u_6^c$, $u_3 u_5^c$, $u_5 u_3^c$, $u_6 u_2^c$ and $u_7 u_7^c$. Since the light spectrum contains $u_{3,4,7}$ and $u_{3,4,7}^c$, the only possible mass term is $H_2 u_7 u_7^c$. The other two up-quark masses can only be generated radiatively, if at all.

Assuming a non-zero vev for H_2 , it is easy to check whether any neutrino masses are generated. The couplings $H_2 L_4 \nu_{6,7}^c$ would appear to generate one neutrino mass of order m_u , which is not suppressed by a see-saw mechanism due to the absence of large Majorana mass terms for the ν^c . None of the quartic superpotential terms in Table 1b can generate such a mass term with the pattern of vev's assumed in this $1(\overline{1}, \overline{2})$ model. It should be noted, however, that the interpretation of the $H_2 L_4 \nu_{6,7}^c$ couplings is not unambiguous, since, as we see in the next subsection, H_2 mixes with the L_4/H_4^c state, which should perhaps be considered a "Higgs" rather than a "lepton" doublet.

Down quark masses can in principle be generated by either H^c or L if either of these acquires a vev. Because all the light charge $+1/3$ "antiquark" states in the $1(\overline{1}, \overline{2})$ model contain combinations of both d^c and D^c states, we may consider either possibility. In the $1(\overline{1}, \overline{2})$ model, the H_2 mixes with both the H_4^c/L_4 and the H_9^c/L_9 states, as discussed in the next subsection, both of which are therefore expected to have non-zero vev's. The Yukawa couplings of Table 1a link (q_4+q_3) to the third state in (14) via H_4^c/L_4 , q_7 to the first state in (14) via H_4^c/L_4 , and (q_4+q_7) to the third state in (14) via H_9^c/L_9 . Since these couplings involve only two of the three charge $+1/3$ "antiquarks" in (14), there are only two massive charge $-1/3$ quarks. Once again, the third charge $-1/3$ quark could only get its mass from radiative corrections, if at all. It should moreover be noted that in most of the other models, even if we got three up-quark masses and three down-quark masses, we would still not be out of the wood, because there are additional quarks and mirror quarks which cannot be massless.

Looking down the first two columns of Table 5, we see that the only model with a full complement of quark masses is $9(\overline{1}, \overline{2})$, whilst the $4(\overline{1}, \overline{2})$ model only fails for the up-quark masses. In the third column, we see that all the models except $7(\overline{1}, \overline{2})$ have insufficient charged lepton masses. The only model which satisfactorily avoids large neutrino masses is seen from the fourth column to be $9(\overline{1}, \overline{2})$. Thus we conclude that most flat models do not have satisfactory fermion masses.

4.2 - HH^C mixing

In conventional supersymmetric phenomenology, the Yukawa coupling of the heaviest quark, presumably the t , drives negative the squared mass of the Higgs field which gives its mass, and thereby generates a vev, namely some $\langle H \rangle \neq 0$. In order to give masses to the charge $-1/3$ quarks, we also need some $\langle H^C/L \rangle \neq 0$, and this occurs naturally if there is a superpotential term mixing HH^C (or HL). Such a mixing term also avoids problems with an axion-like particle. In many models, HH^C mixing arises through a term $\lambda HH^C N$ in the superpotential. If λ (or some other Yukawa coupling of the N field) is sufficiently large, the renormalization group equations show that the squared mass for N will be driven negative, implying $\langle N \rangle \neq 0$. In the present set of models, no such term is possible, since no N fields appear in the light spectrum. It may, however, be possible to have HH^C mixing through higher order terms in the potential such as $HH^C N(N\bar{N})^m/M_C^{2m}$, if such a term exists and $\langle N \rangle^{m+1} \langle \bar{N} \rangle^m / M_C^{2m} \lesssim 1$ TeV when $\langle N \rangle, \langle \bar{N} \rangle \sim m_I$. Analogous terms could also give rise to HL mixing, which carries with it the danger of lepton number violation. With our estimate of $M_I \gtrsim 10^{16}$ GeV, we would judge that HH^C mixing terms with $n \equiv 2m+1 = 5 \pm 1$ are acceptable, whereas terms with $n < 3$ are likely to be too large, and terms with $n > 7$ may be too small. As for HL mixing, terms with $n > 7$ are presumably small enough to be safe, whereas $n = 5 \pm 1$ is marginally acceptable, whilst terms with $n < 3$ would be unacceptable.

In the $1(\bar{1}, \bar{2})$ model, because H_5 does not couple to quarks, the a priori possible and interesting HH^C mixing terms are $H_2 N_1 (N_1 \bar{N}_{1,2})^m \times (H_{2,4,9}^C)$. We find that $H_2 H_2^C$ and $H_2 H_9^C$ mixing are incompatible with the discrete symmetries for any value of n . However, $H_2 H_4^C$ mixing is possible for $n = 5$, and hence could have a phenomenologically acceptable order of magnitude. There is also $H_5 H_5^C$ mixing, but this is not of phenomenological interest for the reason discussed in the previous subsection. In addition, there is $H_2 L_9$ mixing in the same order $n = 5$, which is phenomenologically equivalent to $H_2 H_9^C$ mixing, since one of the light states in (17) is a mixture of L_9 and H_9^C . The only other mixing term in low order is $H_5 L_5$ for $n = 5$, which is marginally acceptable.

The situations in other $n(\bar{1}, \bar{2})$ models are exhibited in Table 5. The $2(\bar{1}, \bar{2})$, $4(\bar{1}, \bar{2})$, $5(\bar{1}, \bar{2})$ and $9(\bar{1}, \bar{2})$ models have acceptable HH^C mixing in order $n = 5$, whilst the $6(\bar{1}, \bar{2})$ and $7(\bar{1}, \bar{2})$ models have HH^C mixing in order $n = 3$, which is problematic. All models have acceptable HL mixing appearing in order $n \geq 5$.

4.3 - Flavour-changing neutral currents (FCNC)

These can be generated by Higgs, lepton or quark exchange, if the Yukawa couplings are not flavour-diagonal. The most familiar case is that of Higgs exchange, where it is well known that to avoid FCNC, all quarks and leptons with the same charge must couple to (and acquire their masses from) the same Higgs field²⁵⁾. This condition is automatically respected in the standard model, which only has one Higgs doublet, and in the minimal supersymmetric standard model, which only has one Higgs doublet $H(H^C)$ of hypercharge $Y = +\frac{1}{2}(-\frac{1}{2})$ giving masses to all the up quarks (down quarks and charged leptons) respectively. A generic superstring model may have many Higgses H and H^C in the light spectrum, all of which may couple to light qq^C or ll^C pairs - in principle there could be as many as 9 in the class of models discussed here. In practice, many of the H and H^C states acquire large masses, and therefore cause no problem, but most of the models discussed in Section 3 have more than one light H and/or H^C doublet. Analogous problems can arise from the lepton-quark-antiquark couplings which are generic features of superstring models. This is because slepton exchanges between quarks and squark exchanges between quark/lepton pairs can both lead to flavour non-conserving interactions. This is another reflection of the interchangeability between H^C and L . Finally, three-quark Yukawa couplings and squark exchange can also lead to flavour-violation if the couplings are not all simultaneously diagonalizable.

Table 5 includes columns for potential flavour non-conservation induced by Hqu^C , H^Cqd^C , H^CLe^C , LqD^C , e^CDu^C , Dqq and $D^Cu^Cd^C$ couplings. It should be noted that none of these problems is necessarily disastrous, since even if the Yukawa couplings are not flavour-conserving, they may be sufficiently small that the magnitudes of FCNC effects are phenomenologically acceptable. Phenomenological upper bounds on these couplings are derived and tabulated in Ref. 12).

We now discuss in some detail the potential FCNC problems in the $1(\overline{1}, 2)$ model. There is no problem with H exchange, because there is only one Hqu^C coupling. In the case of H^C exchanges between quarks, potential problems arise because the light q and d^C states couple to H_2^C , H_4^C , H_5^C and H_9^C . This problem is compounded by the presence of light q and D^C couplings to L_2 , L_4 , L_5 and L_9 . As for H^C exchanges between leptons, there are more possible problems because both H_2^C and H_5^C couple to light L doublets and e^C fields. Finally, we note since there are no light Dqq couplings, there can be no flavour-changing problems with scalar D exchanges, whilst the presence of many $D^Cu^Cd^C$ couplings threatens severe problems with scalar D^C exchanges.

Looking down Table 5, we see that the only model with a clean bill of health for FCNC is the $2(\overline{1}, \overline{2})$ model. Problems are present everywhere else except for H exchanges in the $6(\overline{1}, \overline{2})$ and $7(\overline{1}, \overline{2})$ models, for H^c exchanges between leptons in the $7(\overline{1}, \overline{2})$ model, and for D exchanges in the $9(\overline{1}, \overline{2})$ model.

4.4 - Baryon- and lepton-number-violating interactions

In the standard model, every fermion has either baryon number $B \neq 0$, or lepton number $L \neq 0$, or both, whilst every boson has $B = L = 0$. Hence any model involving just standard model fields which has $\Delta B \neq 0$ must also have $\Delta L \neq 0$, and vice versa. This is no longer true if one augments the standard model spectrum. For example, if there is a stable supersymmetric fermion, say the photino $\tilde{\gamma}$ carrying R-parity, then one can have $\Delta B \neq 0$, $\Delta R \neq 0$, $\Delta L = 0$ and $\Delta B = 0$, $\Delta R \neq 0$, $\Delta L \neq 0$ transitions as well as the $\Delta B \neq 0$, $\Delta L \neq 0$, $\Delta R = 0$ transitions which only involve standard model particles. Here we only worry about contemporary problems, so will only give a full catalogue of problems with $\Delta B \neq 0$, $\Delta L \neq 0$ transitions that could result in baryon decay. Since $m_\tau > m_{p,n} > m_{\mu,e,\nu}$, we will not claim that a model has a problem unless it has $\Delta B, \Delta L \neq 0$ interactions involving neutrinos and/or at least two distinct charged lepton species. In general, such transitions could be mediated either by scalar D^c particles - via $D^c u^c d^c$ and $D^c Lq$ interactions - or by scalar D particles - via Dqq and $De^c u^c$ or $D\nu^c d^c$ interactions. In view of the d^c/D^c and H^c/L mixing which is a general feature of the light particle spectra in Table 4, one should also consider the corresponding generalizations of these exchanges.

In the $1(\overline{1}, \overline{2})$ model there are no Dqq interactions among the light states, so scalar D exchange is no problem. However, there are a large number of $D^c u^c d^c$ and $D^c Lq$ interactions, involving all the light quarks, antiquarks and lepton doublets in the spectrum (14), (15) and (17). Therefore we conclude there is a problem with baryon decay.

Looking down the corresponding columns of Table 5, we see that there are ΔB , $\Delta L \neq 0$ problems everywhere except for the $2(\overline{1}, \overline{2})$ and $9(\overline{1}, \overline{2})$ models. Although we do not discuss them in detail, we also note in passing that all the models have problems with $\Delta B = 0$, $\Delta L \neq 0$, $\Delta R \neq 0$ interactions.

5. - SUMMARY AND CONCLUSIONS

The answer to the question raised in the title of this paper is at best "maybe". In order to obtain the standard model from Calabi-Yau compactification of the heterotic superstring with a rank-6 gauge group in four dimensions, one must generate an intermediate mass scale m_I . This can only be generated by a combination of vev's formed along some flat direction in field space. We have argued that couplings to singlet fields can be ignored in looking for flat directions, of which there are very many in realistic models. Generating a large intermediate scale by some variant of the conventional radiative correction scenario is non-trivial. Even if this can be done, it is difficult to find a consistent value of m_I which respects all the known constraints of baryon stability, renormalization of the gauge couplings, and the absence of excessive entropy generation when the intermediate scale phase transition occurs at $T = O(m_W)$. Even if one can sidestep all these difficulties, one must address the problems raised by the existence of many different flat directions. Why do different parts of the Universe not choose different flat directions, and if all parts of the Universe do live in the same vacuum, how is this chosen? Moreover, any flat direction requires large vev's both for some v^c field(s) and for some N field(s), and can therefore be expected to give a low-energy theory which violates maximally baryon and lepton number. We have looked at a considerable number of specific flat directions, and confirmed that this is indeed the case. We also find that the generic low-energy theory inspired by the superstring possesses many other phenomenological defects. Electroweak gauge symmetry breaking cannot generate masses for all the quarks and/or charged leptons. However, it does generate a Dirac mass for one or more neutrinos, which has no apparent reason to be much smaller than m_q or m_l . The Higgs mixing necessary for a realistic electroweak vacuum may not occur with the appropriate magnitude, and if it does, Higgs-lepton mixing may also occur. Other lepton-number-violating couplings are also generic, as is baryon number violation. Finally, there are often flavour-changing neutral interactions due to the exchanges of multiple Higgses and/or other particles weighing $O(m_W)$.

What is the status of the one specific proposed⁴⁾ model which avoids all these latter difficulties? It suffers from the general difficulty of finding a consistent value of the intermediate scale. There is no a priori reason apparent why the flat direction used in this model should be dynamically preferred over all the others. Finally, some desirable features of the model hinge upon supplementary assumptions about the formation of additional vev's $< m_I$.

Our comments should not be construed as condemning the model of Ref. 4), which exemplifies a very interesting approach with many advantages over other superstring-inspired approaches. Rather, our arguments indicate how special the model is, and how hard one must work to create it. A positive response would be to rise to the challenges of demonstrating an a priori motivation for this flat direction and of avoiding the general difficulties with intermediate scales.

Alternatively, one could duck this challenge and extract the standard model from the superstring in some other way. One possibility would be to postulate²⁶⁾ that some field(s) acquire(s) vev's of order m_p without invoking the flat direction mechanism of Section 2. Another strategy for doing this would be to start from a Calabi-Yau compactification of the heterotic superstring in which the Hosotani mechanism breaks to a rank-5 gauge group in four dimensions. In this case one needs no intermediate mass scale, and the low-energy theory does not violate lepton number independently of baryon number violation. Problems remaining to be investigated include baryon number-violating couplings, quarks, lepton and neutrino masses, and flavour-changing neutral interactions. These are now being studied²⁷⁾ in specific rank-5 models. Another strategy for extracting the standard model would be to work directly with superstring theories formulated directly in four dimensions¹⁰⁾. Presumably, if this formulation was sufficiently general it would include all the compactifications of the heterotic superstring previously considered. However, it could also include other possibilities, maybe even including the standard model $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge group without the need for any intermediate scale of gauge symmetry breaking. Thus, even if the approach to the standard model through a rank-6 Calabi-Yau compactification studied in Ref. 4) and here does not work out, there may be other ways in which the superstring can inspire the standard model.

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TABLE 1

a) Cubic Superpotential Terms

27^3

$$\lambda^3 \left\{ \begin{array}{l} \lambda_1^3 + \lambda_2^3 \\ (\lambda_1 \lambda_3 + \lambda_2 \lambda_4) \lambda_6 \\ (\lambda_1 \lambda_3 - \lambda_2 \lambda_4) \lambda_7 \\ \lambda_1 \lambda_4 \lambda_9 + \lambda_2 \lambda_3 \lambda_8 \\ \lambda_3^3 + \lambda_4^3 \\ (\lambda_5 + \lambda_6/2)^2 \lambda_6 \\ \lambda_6 \lambda_7^2 \\ \lambda_6 \lambda_8 \lambda_9 \\ \lambda_8^3 + \lambda_9^3 \end{array} \right.$$

$$q^3 \left\{ \begin{array}{l} q_1^3 \\ q_2^3 \\ q_3^3 \\ q_4^3 + q_5^3 \\ q_6^3 + q_7^3 \\ q_1 q_2 q_3 \\ q_1 q_4 q_5 \\ q_2 q_6 q_7 \\ q_3 (q_4 q_6 + q_5 q_7) \end{array} \right.$$

$$Q^3 \left\{ q \rightarrow Q \right.$$

$$\lambda q Q \left\{ \begin{array}{l} \lambda_1 q_2 Q_7 + \lambda_2 q_2 Q_6 \\ \lambda_1 q_3 Q_4 + \lambda_2 q_3 Q_5 \\ \lambda_1 q_4 Q_3 + \lambda_2 q_5 Q_3 \\ \lambda_1 q_6 Q_6 + \lambda_2 q_7 Q_7 \\ \lambda_1 q_7 Q_2 + \lambda_2 q_6 Q_2 \\ \lambda_3 q_1 Q_5 + \lambda_4 q_1 Q_4 \\ \lambda_3 q_5 Q_1 + \lambda_4 q_4 Q_1 \\ \lambda_3 q_3 Q_6 + \lambda_4 q_3 Q_7 \\ \lambda_3 q_6 Q_3 + \lambda_4 q_7 Q_3 \\ \lambda_3 q_4 Q_4 + \lambda_4 q_5 Q_5 \\ (\lambda_5 + \lambda_6/2) q_1 Q_2 \\ -(\lambda_5 + \lambda_6/2) q_2 Q_1 \\ \lambda_6 q_3 Q_3 \\ (\lambda_5 + \lambda_6/2) (q_4 Q_6 + q_5 Q_7) \\ -(\lambda_5 + \lambda_6/2) (q_6 Q_4 + q_7 Q_5) \\ \lambda_7 (q_4 Q_6 - q_5 Q_7) \\ \lambda_7 (q_6 Q_4 - q_7 Q_5) \\ \lambda_8 q_1 Q_6 + \lambda_9 q_1 Q_7 \\ \lambda_8 q_6 Q_1 + \lambda_9 q_7 Q_1 \\ \lambda_8 q_2 Q_5 + \lambda_9 q_2 Q_4 \\ \lambda_8 q_5 Q_2 + \lambda_9 q_4 Q_2 \\ \lambda_8 q_4 Q_7 + \lambda_9 q_5 Q_6 \\ \lambda_8 q_7 Q_4 + \lambda_9 q_6 Q_5 \end{array} \right.$$

$\overline{27}^3$

$$\bar{\lambda}^3 \left\{ \begin{array}{l} (\bar{\lambda}_1, \bar{\lambda}_2)^3 \\ (\bar{\lambda}_1, \bar{\lambda}_2) (\bar{\lambda}_3, \bar{\lambda}_4) (\bar{\lambda}_5, \bar{\lambda}_6) \end{array} \right.$$

$$\bar{\lambda} \bar{q} \bar{Q} \left\{ \begin{array}{l} (\bar{\lambda}_1, \bar{\lambda}_2) (\bar{q}_1 \bar{Q}_2 + \bar{q}_2 \bar{Q}_1) \\ (\bar{\lambda}_1, \bar{\lambda}_2) (\bar{q}_1 \bar{Q}_4 + \bar{q}_2 \bar{Q}_3) \\ (\bar{\lambda}_1, \bar{\lambda}_2) (\bar{q}_3 \bar{Q}_2 + \bar{q}_4 \bar{Q}_1) \\ (\bar{\lambda}_1, \bar{\lambda}_2) (\bar{q}_3 \bar{Q}_4 + \bar{q}_4 \bar{Q}_3) \end{array} \right.$$

$$\bar{q}^3 \left\{ \begin{array}{l} (\bar{q}_1, \bar{q}_3)^3 \\ (\bar{q}_2, \bar{q}_4)^3 \end{array} \right.$$

$$\bar{Q}^3 \left\{ \bar{q} \rightarrow \bar{Q} \right.$$

TABLE 1

b) Quartic Superpotential Terms

$$\begin{aligned}
 & \lambda^2 \bar{\lambda}^2 \left\{ \begin{aligned} P_1 &= (\lambda_1 \lambda_2) (\bar{\lambda}_3, \bar{\lambda}_4)^2 \\ P_2 &= (\lambda_1 \lambda_3 + \lambda_2 \lambda_4) [(\bar{\lambda}_1, \bar{\lambda}_2)^2, \bar{\lambda}_3, \bar{\lambda}_4, \bar{\lambda}_5, \bar{\lambda}_6] \\ P_3 &= (\lambda_1 \lambda_8 + \lambda_2 \lambda_9) (\bar{\lambda}_5, \bar{\lambda}_6)^2 \\ P_5 &= (\lambda_3 \lambda_4) (\bar{\lambda}_5, \bar{\lambda}_6)^2 \\ P_6 &= (\lambda_3 \lambda_9 + \lambda_4 \lambda_8) (\bar{\lambda}_3, \bar{\lambda}_4)^2 \\ P_8 &= [(\lambda_5 + \lambda_6/2)^2, \lambda_7^2] [(\bar{\lambda}_1, \bar{\lambda}_2)^2, \bar{\lambda}_3, \bar{\lambda}_4, \bar{\lambda}_5, \bar{\lambda}_6] \\ P_9 &= (\lambda_8 \lambda_9) [(\bar{\lambda}_1, \bar{\lambda}_2)^2, \bar{\lambda}_3, \bar{\lambda}_4, \bar{\lambda}_5, \bar{\lambda}_6] \end{aligned} \right. \\
 & q^2 \bar{q}^2 \left\{ \begin{aligned} & (q_1 q_2) (\bar{q}_1, \bar{q}_3, \bar{q}_2, \bar{q}_4) \\ & (q_1 q_6) (\bar{q}_2, \bar{q}_4)^2 \\ & (q_1 q_7) (\bar{q}_1, \bar{q}_3)^2 \\ & (q_2 q_4) (\bar{q}_1, \bar{q}_3)^2 \\ & (q_2 q_5) (\bar{q}_2, \bar{q}_4)^2 \\ & (q_3^2) (\bar{q}_1, \bar{q}_3, \bar{q}_2, \bar{q}_4) \\ & (q_4 q_6) (\bar{q}_1, \bar{q}_3, \bar{q}_2, \bar{q}_4) \\ & (q_4 q_7) (\bar{q}_2, \bar{q}_4)^2 \\ & (q_5 q_6) (\bar{q}_1, \bar{q}_3)^2 \\ & (q_5 q_7) (\bar{q}_1, \bar{q}_3, \bar{q}_2, \bar{q}_4) \end{aligned} \right. \\
 & \lambda \bar{\lambda} q \bar{q} \left\{ \begin{aligned} & (\lambda_1 \bar{\lambda}_1, 2) (q_1 \bar{q}_1, 3 + q_5 \bar{q}_2, 4) \\ & (\lambda_1 \bar{\lambda}_3, 4) (q_3 \bar{q}_1, 3) \\ & (\lambda_1 \bar{\lambda}_5, 6) (q_2 \bar{q}_1, 3 + q_6 \bar{q}_2, 4) \\ & (\lambda_2 \bar{\lambda}_1, 2) (q_1 \bar{q}_2, 4 + q_4 \bar{q}_1, 3) \\ & (\lambda_2 \bar{\lambda}_3, 4) (q_3 \bar{q}_2, 4) \\ & (\lambda_2 \bar{\lambda}_5, 6) (q_2 \bar{q}_2, 4 + q_7 \bar{q}_1, 3) \\ & (\lambda_3 \bar{\lambda}_1, 2) (q_1 \bar{q}_2, 4 + q_4 \bar{q}_1, 3) \\ & (\lambda_3 \bar{\lambda}_3, 4) (q_1 \bar{q}_2, 4 + q_4 \bar{q}_1, 3) \\ & (\lambda_3 \bar{\lambda}_5, 6) (q_3 \bar{q}_2, 4) \\ & (\lambda_4 \bar{\lambda}_1, 2) (q_2 \bar{q}_1, 3 + q_6 \bar{q}_2, 4) \\ & (\lambda_4 \bar{\lambda}_3, 4) (q_1 \bar{q}_1, 3 + q_5 \bar{q}_2, 4) \\ & (\lambda_4 \bar{\lambda}_5, 6) (q_3 \bar{q}_1, 3) \\ & (\lambda_6 \bar{\lambda}_3, 4) (q_6 \bar{q}_1, 3 - q_7 \bar{q}_2, 4) \\ & (\lambda_6 \bar{\lambda}_5, 6) (q_4 \bar{q}_2, 4 - q_5 \bar{q}_1, 3) \\ & (\lambda_7 \bar{\lambda}_3, 4) (q_6 \bar{q}_1, 3 + q_7 \bar{q}_2, 4) \\ & (\lambda_7 \bar{\lambda}_5, 6) (q_6 \bar{q}_2, 4 + q_7 \bar{q}_1, 3) \\ & (\lambda_8 \bar{\lambda}_1, 2) (q_3 \bar{q}_2, 4) \\ & (\lambda_8 \bar{\lambda}_3, 4) (q_2 \bar{q}_2, 4 + q_7 \bar{q}_1, 3) \\ & (\lambda_8 \bar{\lambda}_5, 6) (q_1 \bar{q}_2, 4 + q_4 \bar{q}_1, 3) \\ & (\lambda_9 \bar{\lambda}_1, 2) (q_3 \bar{q}_1, 3) \\ & (\lambda_9 \bar{\lambda}_3, 4) (q_2 \bar{q}_1, 3 + q_6 \bar{q}_2, 4) \\ & (\lambda_9 \bar{\lambda}_5, 6) (q_1 \bar{q}_1, 3 + q_5 \bar{q}_2, 4) \end{aligned} \right. \\
 & \left. \lambda \bar{\lambda} q \bar{q} \right\} \left\{ \begin{aligned} & q, \bar{q} \rightarrow Q, \bar{Q} \\ & q^2 \bar{q}^2 \end{aligned} \right.
 \end{aligned}$$

TABLE 2

Terms making possibly non-flat contributions to V

$\lambda_i \backslash \bar{\lambda}_j$	$\bar{1}$	$\bar{2}$	$\bar{3}$	$\bar{4}$	$\bar{5}$	$\bar{6}$
1	P_2	P_2	P_1, P_2^2	P_1, P_2^2	P_3^3	P_3^3
2	P_2	P_2	P_1, P_2^2	P_1, P_2^2	P_3^3	P_3^3
3	P_2^3	P_2^3	P_2, P_6^2	P_2, P_6^2	P_2, P_3, P_5^2	P_2, P_3, P_5^2
4	P_2	P_2	P_2, P_6^2	P_2, P_6^2	P_3, P_5^2	P_3, P_5^2
5	P_8	P_8	P_2, P_8^2	P_2, P_8^2	P_3, P_8	P_3, P_8
6	P_8	P_8	P_2, P_8^2	P_2, P_8^2	P_3, P_8	P_3, P_8
7	P_8	P_8	P_2, P_8^2	P_2, P_8^2	P_3, P_8	P_3, P_8
8	P_3, P_9	P_3, P_9	P_2, P_3, P_6, P_9^2	P_2, P_3, P_6, P_9^2	P_3^3	P_3^3
9	P_9	P_9	P_2, P_6, P_9	P_2, P_6, P_9	P_3^3	P_3^3

The superstripts indicate the number of possible contributions provided by the corresponding term in Table 1(b).

TABLE 3

a) Down Quark Mass Matrix in $1(\overline{1}, \overline{2})$ Model

		d^c							D^c							$\overline{D}$			
		1	2	3	4	5	6	7	1	2	3	4	5	6	7	1	2	3	4
D	1															Q		Q	
	2							C							C	Q		Q	
	3				C						C								
	4			C							C								
	5																Q		Q
	6						C							C		Q		Q	
	7		C						C										
$\overline{d}^c$	1	Q	Q																
	2					Q	Q												
	3	Q	Q																
	4					Q	Q												
$\overline{D}^c$	1								Q	Q						C		C	
	2											Q	Q			C		C	
	3								Q	Q							C		C
	4												Q	Q		C		C	

TABLE 3

b) Lepton/Higgs Doublet Mass Matrix in $1(\bar{1},2)$ Model

		H									$\overline{L}$						$\overline{H}^c$						$\tilde{G}$	
		1	2	3	4	5	6	7	8	9	1	2	3	4	5	6	1	2	3	4	5	6		
L	1	C																						h
	2																							
	3						C	C			Q ₂	Q ₂	Q ₂	Q ₂										
	4									C														
	5																							
	6			C																				
	7			C																				
	8													Q ₃	Q ₃									
	9				C																			
H^c	1	C																					h	
	2																							
	3						C	C									Q ₂	Q ₂	Q ₂	Q ₂				
	4									C														
	5																							
	6			C																				
	7			C																				
	8																				Q ₃	Q ₃		
	9				C																			
$\overline{H}$	1			Q ₂									C	C			C	C						
	2			Q ₂									C	C			C	C						
	3			Q ₂							C	C								C	C			
	4			Q ₂							C	C								C	C			
	5									Q ₃									C	C				
	6									Q ₃									C	C				
$\tilde{G}$											h h						h h							

TABLE 4

a) Light Lepton States

	$L/H^c/\bar{H}$	$H/\bar{L}/\bar{H}^c$	e^c	$\bar{e}^c$	ν^c
1	$L_{2,5}$ $H_{2,5}^c$ $1/(L_4, H_4^c), 1/(L_9, H_9^c)$	$H_{2,5}$ $\bar{L}_{5,6}$	$e_{2,4,5,6,7,9}^c$	$2/(\bar{e}_{1,2,3,4}^c)$ $1/(\tilde{G}_{1,2})$	$\nu_{5,6,7}^c$
2	$L_5, 1/(L_2, H_1^c)$ $H_{5,9}^c$ $1/(L_4, 6, 7, H_8^c, \bar{H}_{3,4}, \tilde{G})$	H_5 $1/(\bar{L}_{5,6})$	$e_{5,6,7,9}^c$	$1/(\bar{e}_{3,4}^c)$	$\nu_{5,6,7}^c$
4	$L_5, 1/(L_1, H_9^c)$ $H_{5,1}^c/(L_9, H_4^c)$ $2/(L_2, 6, 7, H_8^c, \bar{H}_{3,4}, \tilde{G})$	H_5 $1/(H_{2,6,7}, \bar{H}_{1,2}^c, \bar{L}_{3,4})$ $1/(\bar{L}_{5,6})$	$e_{5,6,7,9}^c$	$1/(\bar{e}_{3,4}^c)$	$\nu_{5,6,7}^c$
5	L_2 $H_{2,4,7,9}^c$ $2/(L_{5,6,7}, H_{1,8}^c, \bar{H}_{3,4}, \tilde{G})$	H_2 $2/(H_{5,6,7}, \bar{L}_{3,4}, \bar{H}_{1,2}^c)$ $1/(\bar{L}_{5,6})$	$e_{2,4,7,9}^c$	$1/(\bar{e}_{3,4}^c)$	ν_7^c
6	$L_2, 1/(L_9, H_2^c)$ $1/(L_1, H_{3,9}^c, \bar{H}_{5,6})$ $2/(L_4, 5, 6, 7, H_{1,8}^c, \bar{H}_{3,4}, \tilde{G})$	$1/(H_{5,6,7}, \bar{L}_{3,4}, \bar{H}_{1,2}^c)$ $1/(\bar{L}_{5,6})$	$e_{2,4,7,9}^c$	$1/(\bar{e}_{3,4}^c)$	ν_7^c
7	$L_{2,5}, H_{5,9}^c$ $1/(L_9, H_2^c)$ $1/(L_6, 7, H_{1,8}^c, \bar{H}_{3,4}, \tilde{G})$	H_5 $1/(\bar{L}_{5,6})$ $1/(H_6, 7, \bar{L}_{3,9}, \bar{H}_{1,2}^c)$	$e_{2,4,5,6,9}^c$	$1/(\bar{e}_{3,4}^c)$ $1/(\bar{e}_{1,2,3}^c, \tilde{G})$	$\nu_{5,6}^c$
9	$L_{2,5}$ $H_{2,5,7}^c$ $1/(L_1, H_4^c), 1/(L_4, H_9^c)$	$H_{2,5}$ $2/(H_{3,6,7,8},$ $\bar{L}_{1,2,3,4,5,6},$ $\bar{H}_{1,2,3,4,5,6}^c)$	$e_{2,4,5,6,7}^c$	$1/(\bar{e}_{3,4}^c)$ $1/(\bar{e}_{1,2,5,6}^c)$	$\nu_{5,6,7}^c$

TABLE 4

b) Light Quark States

	q	$\bar{q}$	$d^c/D^c/\bar{D}$	$D/\bar{d}^c/\bar{D}^c$	u^c	$\bar{u}^c$
1	$q_{3,4,7}$	-	$1/(d_3^c, D_3^c), 1/(d_4^c, D_4^c)$ $1/(d_{1,2,5,6,7}^c,$ $D_{1,2,5,6,7}^c, \bar{D}_{1,2,3,4})$	-	$u_{3,4,7}^c$	-
2	$q_{3,5,7}$	-	d_5^c $2/(d_{3,7}^c, D_{1,6}^c, \bar{D}_{1,3})$	-	$u_{3,5,7}^c$	-
4	$q_{1,3,4,5,7}$	$1/(\bar{q}_{1,3})$ $1/(\bar{q}_{2,4})$	$d_{1,5}^c$ $1/(d_3^c, D_1^c), 1/(d_4^c, D_7^c)$ $1/(d_7^c, D_6^c, \bar{D}_{1,3})$	$1/(D_7, \bar{d}_{1,3}^c)$ $1/(D_6, \bar{d}_{2,4}^c, \bar{D}_{1,3}^c)$	$u_{1,3,4,5,7}^c$	$1/(\bar{u}_{1,3}^c)$ $1/(\bar{u}_{2,4}^c)$
5	$q_{1,3,4,5,7}$	$1/(\bar{q}_{1,3})$ $1/(\bar{q}_{2,4})$	$d_{1,5}^c, D_3^c$ $2/(d_{3,7}^c, D_{1,6}^c, \bar{D}_{1,3})$	$1/(D_7, \bar{d}_{1,3}^c)$ $1/(D_{1,6}, \bar{d}_{2,4}^c, \bar{D}_{1,3}^c)$	$u_{1,3,4,5,7}^c$	$1/(\bar{u}_{1,3}^c)$ $1/(\bar{u}_{2,4}^c)$
6	$q_{1,3,4,5,7}$	$1/(\bar{q}_{1,3})$ $1/(\bar{q}_{2,4})$	$d_{1,5}^c$ $1/(d_4^c, D_3^c)$ $2/(d_{3,7}^c, D_{1,6}^c, \bar{D}_{1,3})$	$1/(D_7, \bar{d}_{1,3}^c)$ $1/(D_6, \bar{d}_{2,4}^c, \bar{D}_{1,3}^c)$	$u_{1,3,4,5,7}^c$	$1/(\bar{u}_{1,3}^c)$ $1/(\bar{u}_{2,4}^c)$
7	$q_{1,3,4,5,7}$	$1/(\bar{q}_{1,3})$ $1/(\bar{q}_{2,4})$	$d_{1,5}^c$ $D_{1,3}^c$ $1/(d_{3,7}^c, D_6^c, \bar{D}_{1,3})$	D_1 $1/(D_7, \bar{d}_{1,3}^c)$	$u_{1,3,4,5,7}^c$	$1/(\bar{u}_{1,3}^c)$ $1/(\bar{u}_{2,4}^c)$
9	$q_{1,4,5,7}$	$1/(\bar{q}_{2,4})$	$1/d_{1,5}^c$ $1/(\bar{d}_{4,7}^c, D_{4,6}^c, \bar{D}_{1,3})$	-	$u_{1,4,5,7}^c$	$1/(\bar{u}_{2,4}^c)$

TABLE 5
Possible problems in $n(1,2)$ flat directions

n	$3m_u$ + mirrors	$3m_d$ + mirrors	$3m_\ell$ + mirrors	$m_\nu=0$	HH^c mixing	no HL mixing	no $\Delta B \neq 0$	no $\Delta L \neq 0$	Absence of FCNC due to				
									Hq_u^c	$H^c q_d^c$	$H^c \ell_e^c$	Dq_q	$D^c u^c d^c$
1	X(1)	X(2)	X(1)	X(1)	✓ ⁺	✓ ^x	X	X	✓	X	X	✓	X
2	X(2)	X(2)	X(2)	X(2)	✓ ⁺	✓ ^x	✓	X	✓	✓	✓	✓	✓
4	X(3) [*]	✓	X(2)	X(3)	✓ ⁺	✓ ^x	X	X	X	X	X	X	X
5	X(3) [*]	X(4) [*]	X(1)	X(1)	✓ ⁺	✓	X	X	X	X	X	X	X
6	X(3) [*]	X(1) [*]	X(3) [*]	X(1)	X ^x	✓	X	X	✓	X	X	X	X
7	X(3) [*]	X(2) [*]	✓	X(2)	X ^x	✓	X	X	✓	X	✓	X	X
9	✓	✓	X(2)	✓	✓ ⁺	✓ ^x	✓	X	X	X	X	✓	X

The numbers in parentheses indicate the number of massive states.

*: not all mirror particles acquire masses

+ : mixing in order $< > 5$.

x : mixing in order $< > 3$.

Problems are denoted by X, their absence by ✓.

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