GMSB at a stable vacuum and MSSM without exotics from heterotic string

Jihn E. Kim

Department of Physics and Astronomy and Center for Theoretical Physics, Seoul National University, Seoul 151-747, Korea

Abstract

We show that it is possible to introduce the confining hidden sector gauge group $SU(5)'$ with the chiral matter $10'_0$ plus $5'_0$, which are neutral under the standard model gauge group, toward a gauge mediated supersymmetry breaking (GMSB) in a $\mathbb{Z}_{12-1}$ orbifold compactification of $E_8 \times E'_8$ heterotic string. Three families of MSSM result without exotics. We also find a desirable matter parity $P$ (or $R$-parity) assignment. We note that this model contains the spectrum of the Lee-Weinberg model which has a nice solution of the $\mu$ problem.

PACS numbers: 11.25.Mj, 11.25.Wx, 12.60.Jv

Keywords: Gauge mediation, DSB, Stable vacuum, Orbifold compactification, Matter parity, $\mu$ solution
I. INTRODUCTION

The supersymmetric (SUSY) extension of the standard model (SM) encounters a few naturalness problems, the SUSY flavor problem \[1\], the little hierarchy problem \[2\], the $\mu$ problem \[3\], etc. The hierarchical magnitude is worst in the $\mu$ problem but here there are nice solutions \[4\]. The little hierarchy problem has weakened the nice feature of the SUSY solution of the gauge hierarchy problem and we hope that it will be understood somehow in the future. On the other hand, the SUSY flavor problem seems to require family independence of the interactions at the GUT scale. The attractive gravity mediation scenario for transmitting SUSY breaking down to the observable sector probably violate the flavor independence of interactions violently. This observation has led to the gauge mediated supersymmetry breaking (GMSB) \[5\]. However, the superstring attempt toward a GMSB model has not been successful phenomenologically, even though the possibility of SUSY breaking spectra was pointed out \[6\].

Recently, dynamical SUSY breaking (DSB) at an unstable minimum at the origin of the field space got quite an interest following Intrilligator, Seiberg and Shih (ISS) \[7, 8, 9\], partly because it has not been successful in deriving a phenomenologically attractive model in the stable vacuum. Among the results on SU($N$), SO($N$) and Sp($2n$) groups, the result is especially simple for SU($N_c$) with $N_f$ flavors, showing an unstable minimum for $N_c + 1 \leq N_f < \frac{3}{2}N_c$. This mechanism is easily applicable to SU($5'$) models with 6 or 7 flavors, which can be realized in string compactifications \[9\]. Nevertheless, it is better to realize a phenomenologically successful SUSY breaking *stable minimum*, not to worry about our stability in a remote future. In this paper, therefore, we look for a GMSB spectrum in the orbifold compactification of the $E_8 \times E_8'$ heterotic string with three families, trying to satisfy all obvious phenomenological requirements.

The well-known DSB models are an SO($10'$) model with 16' or 16' + 10' \[10\], and an SU($5'$) model with 10' + 5' \[11\]. It is known that GMSB with 16' + 10' can be obtained from heterotic string \[12\], but the beta function magnitude is too large (in the negative) so that SO($10'$) confines somewhat above $10^{13}$ GeV against a meaningful GMSB. If the hidden sector gauge group is large, the content of matter representation is usually small and the beta function magnitude (in the negative) turns out to be too large to implement the GMSB scenario. If the confining group is SU(4)' or smaller, it is not known that one can obtain...
a SUSY breaking stable minimum. Thus, SU(5)′ is an attractive choice for the GMSB \cite{6}. To solve the SUSY flavor problem along this line of the GMSB, we require two conditions: relatively low hidden sector confining scale ($\lesssim 10^{12}$ GeV) and appearance of matter spectrum allowing SUSY breaking.

A nice feature of the ISS type model at an unstable vacuum toward model building is that the SUSY breaking can be mediated through dimension-4 superpotential given in\(^1\)

$$W \sim \frac{1}{M} Q\overline{Q}f\bar{f}$$

where $Q$ is a hidden sector quark and $f$ is a messenger. It is possible because the vectorlike representations, for example six or seven ($Q + \overline{Q}$), are present and the $Q\overline{Q}f\bar{f}$ interaction is suppressed by one power of mass parameter. So this mass parameter can be raised up to the GUT scale.

On the other hand, the uncalculable model with $10' + \overline{5}'$ of SU(5)' does not have such a simple singlet direction in terms of chiral fields. For example, the term $\epsilon_{ijklm}10^{ij}10^{kl}10^{mn}\overline{5}_n = 0$ since taking $n = 1$ without generality it is proportional to $\epsilon_{ijklm}10^{ij}10^{kl}10^{m1}\overline{5}_1$ which can be shown to be vanishing using the antisymmetric symbol $\epsilon$. The singlet combination is possible in terms of the chiral gauge field strength, $W'_{\alpha}W'_{\alpha}$. It is pointed out that the $F$-term of this singlet combination can trigger the SUSY breaking to low energy \cite{13},

$$\mathcal{L} = \int d^2\theta \left( \frac{1}{M^2} f\bar{f}W'_{\alpha}W'_{\alpha} + M_f f\bar{f} \right) + \text{h.c.}$$

where the effective parameters of $M$ and $M_f$ can be lower than the GUT scale.

The GMSB problem in string models is very interesting. For example, quite recently but before ISS, it has been reviewed \cite{14}, but the phenomenological requirements toward the minimal supersymmetric standard model (MSSM) have made it difficult to be found in string models. The three family condition works as a strong constraint in the search of the hidden sector representations. If we require the exotics free condition, the possibility reduces dramatically.

In a $\mathbb{Z}_{12-1}$ orbifold compactification, we find a model achieving the GMSB at a stable vacuum together with three families of quarks and leptons without any exotics. Since there is no exotics, it is hoped that the singlet VEVs toward successful Yukawa couplings have much

\(^1\) This form has been considered by many \cite{4}, in particular in \cite{8}.
more freedom, most of which are set at the string scale. We find a successful embedding of
matter parity $P$ and a nice solution of the $\mu$ problem. One unsatisfactory feature is that
$\sin^2 \theta_W$ is not $\frac{3}{8}$. Thus, to fit the weak mixing angle to the observed value, we must assume
intermediate state vectorlike particles. Anyway, another kind of intermediate state particles
is needed also for a successful messenger mass scale.

II. A $\mathbb{Z}_{12-I}$ MODEL

The twist vector in the six dimensional (6d) internal space is

$$\mathbb{Z}_{12-I} \text{ shift : } \phi = \left( \frac{5}{12} \frac{4}{12} \frac{1}{12} \right). \quad (1)$$

We obtain the 4D gauge group by considering massless conditions satisfying $P \cdot V = 0$ and
$P \cdot a_3 = 0$ in the untwisted sector [15]. We embed the discrete action $\mathbb{Z}_{12-I}$ in the $E_8 \times E_8$
space in terms of the shift vector $V$ and the Wilson line $a_3$ as

$$ V = \frac{1}{12}(6 \ 6 \ 6 \ 2 \ 2 \ 2 \ 3 \ 3 \ 3 \ 3 \ 1 \ 1 \ 1)' \quad (2) $$
$$ a_3 = \frac{1}{3}(1 \ 1 \ 2 \ 0 \ 0 \ 0 \ 0 \ 0 \ 0 \ 0 \ 1 \ 1 \ -2)' \quad (3) $$

(a) Gauge group: The 4D gauge groups are obtained by $P^2 = 2$ vectors satisfying $P \cdot V = 0$
and $P \cdot a_3 = 0 \mod \text{integer},$

$$ SU(3)_c \times SU(3)_W \times SU(2)_n \times U(1)_a \times U(1)_b \times U(1)_c $$
$$ \times [SU(5)'/SU(3)' \times U(1)'/2]. \quad (4) $$

The gauge group $SU(3)_W$ will be broken down to $SU(2)_W$ by the vacuum expectation value
(VEV) of $3$ and $\bar{3}$ of $SU(3)_W$. Then, the simple roots of our interest $SU(3)_c$, $SU(2)_W$, and
$SU(2)_n$ are

$$ SU(3)_c : \begin{cases} \alpha_1 = (1 -1 0 0 0 0 0) \\ \alpha_2 = (0 1 1 0 0 0 0) \end{cases} \quad (5) $$
$$ SU(2)_W : \begin{cases} \alpha_1 = (0 0 0 1 -1 0 0) \end{cases} \quad (6) $$
$$ SU(2)_n : \begin{cases} \alpha_1 = (0 0 0 0 0 1 -1) \end{cases} \quad (7) $$
The hypercharge direction is the combination of U(1)s of Eq. (4) and some generators of nonabelian groups

\[ Y = Y_{\text{Abel}} + \frac{1}{\sqrt{3}} W_8 + F_3 - \frac{1}{\sqrt{3}} F_8 = \tilde{Y} + F_3 - \frac{1}{\sqrt{3}} F_8 \]  

(8)

where

\[ Y_{\text{Abel}} = Y_8 + Y_8', \]  

(9)

and \( W_8, F_3, F_8 \) are nonabelian generators of SU(3)_W and SU(3)'. We define \( \tilde{Y} = Y_{\text{Abel}} + \frac{1}{\sqrt{3}} W_8 \) by including the U(1) generators of SU(3)_W and SU(2)_V (by VEVs of scalar fields). \( Y_8 \) and \( Y_8' \) are a linear combination of three U(1) generators in \( E_8 \) and a linear combination of two U(1) generators in \( E'_8 \), respectively. \( W_8 \) is the eighth generator of SU(3)_W, \( \left( \frac{1}{\sqrt{3}}, \frac{1}{\sqrt{3}}, -\frac{1}{\sqrt{3}} \right) \), and \( F_3 \) and \( F_8 \) are the third and the eighth generators of SU(3)', \( \left( \frac{1}{2}, -\frac{1}{2}, 0 \right) \) and \( \left( \frac{1}{2\sqrt{3}}, \frac{1}{2\sqrt{3}}, -\frac{1}{\sqrt{3}} \right) \), respectively. We find that exotics cannot be made vectorlike if we do not include \( Y' \). \( \tilde{Y} \) is defined as

\[ \tilde{Y} = Y_{\text{Abel}} + \frac{1}{\sqrt{3}} W_8 = \left( \frac{1}{6} \frac{1}{6} -\frac{1}{6} 0 0 \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{6} \frac{1}{6} \right) \left( \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{6} \frac{1}{6} \right)' \]  

(10)

We included the SU(3)' generators in \( Y \) of (8) so that there does not appear exotics.

The five U(1) generators of (4) are defined as

\[
\begin{align*}
Q_1 &= (6 6 -6 0 0 0 0 0)(0 0 0 0 0 0 0 0)' \\
Q_2 &= (0 0 0 6 6 6 0 0)(0 0 0 0 0 0 0 0)' \\
Q_3 &= (0 0 0 0 0 2 2)(0 0 0 0 0 0 0 0)' \\
Q_4 &= (0 0 0 0 0 0 0)(4 4 4 4 4 0 0 0)' \\
Q_5 &= (0 0 0 0 0 0 0)(0 0 0 0 0 4 4 4)' 
\end{align*}
\]  

(11)

(b) Matter representations: Now there is a standard method to obtain the massless spectrum in \( Z_{12-I} \) orbifold models. The spectra in the untwisted sectors \( U_1, U_2, \) and \( U_3, \) and twisted sectors, \( T_{10, +}, T_{20, +}, T_{3}, T_{40, +}, T_{50, +}, \) and \( T_6, \) are easily obtained [16]. The representations are denoted as

\[
[\text{SU(3)}_c, \text{SU(2)}_W; \text{SU(5)}', \text{SU(3)}']_{\tilde{Y}},
\]  

(12)

where we already use the broken SU(3)_W and \( \tilde{Y} = Y_{\text{Abel}} + \frac{1}{\sqrt{3}} W_8 \) given in Eq. (10). For obvious cases, we will use the abbreviated notation

\[
(\text{SU(3)}_c, \text{SU(2)}_W)_{\tilde{Y}}.
\]  

(13)
But when SU(3)' triplets or antitriplets are involved, the hypercharge is $\tilde{Y}$. We list all matter fields below,

$$
U_1 : (1, 2)_{1/2}, \ 2 \cdot (1, 2)_{-1/2}, \ 1_1, \ 2 \cdot 1_0 \\
U_2 : (1, 2)_{-1/2}, \ 1_0 \\
U_3 : (1, 2)_{-1/2}, \ 2 \cdot (1, 2)_{1/2}, \ 2 \cdot 1_1 \\
T_{10} : (\bar{3}, 1)_{1/3}, \ (1, 2)_{1/2}, \ 3 \cdot 1_1, \ 2 \cdot 1_0 \\
T_{1-} : (1; \bar{3}, 1, 1, 3')_{1/3}, \ 2 \cdot 1_{-1} \\
T_{20} : (3, 1)_{1/3}, \ (1, 2)_{-1/2}, \ 3 \cdot 1_0 \\
T_{2+} : (1; 10', 1)_0, \ (1; 1, 3')_{1/3}, \ 4 \cdot 1_0 \\
T_3 : 2 \cdot (1; 5', 1)_0, \ 2 \cdot (1; \bar{5}', 1)_0, \\
(2L + 1_L)(1, 2)_{1/2}, \ (1L + 2_R)(1, 2)_{1/2}, \\
(2L + 1_R)1_1, \ 3 \cdot 1_0, \ (6L + 6R) \cdot 1_1 \\
T_{4o} : 3 \cdot (1, 2; 1, \bar{3}')_{1/6}, \ 3 \cdot (1; 1, 3')_{-1/3} \\
T_{4+} : 5 \cdot (1; 1, 3')_{1/3}, \ 2 \cdot (1; 1, 3')_{-2/3} \\
T_4 : 3 \cdot (3, 2)_{1/6}, \ 2 \cdot (\bar{3}, 1)_{-2/3}, \ 5 \cdot (3, 1)_{1/3}, \ 3 \cdot (3, 1)_{-1/3}, \\
5 \cdot (1, 2)_{-1/2}, \ 2 \cdot (1, 2)_{1/2}, \ 2 \cdot 1_1, \ 12 \cdot 1_0, \ 12 \cdot 1_0 \\
T_{70} : (1; \bar{3}', 1)_0, \ (1; 1, 3')_{-2/3} \\
T_{7+} : (\bar{3}, 1)_{-2/3}, \ (3, 1)_{-1/3}, \ 2 \cdot (1, 2)_{-1/2}, \ 1_0, \ 3 \cdot 1_{-1} \\
T_7 : (1; 5', 1)_0, \ (1; 1, \bar{3}')_{-1/3}, \ 2 \cdot 1_1 \\
T_6 : 3 \cdot (1; 5', 1)_0, \ 3 \cdot (1; \bar{5}', 1)_0, \ 2 \cdot (1; 5', 1)_1, \ 2 \cdot (1; \bar{5}', 1)_{-1}
$$

where $1 = (1, 1, 1; 1, 1)$. Breaking SU(3)', we assign

$$
F_3 = (\frac{1}{2}, -\frac{1}{2}, 0), \quad \frac{1}{\sqrt{3}} F_8 = (\frac{1}{6}, \frac{1}{6}, -\frac{1}{3}).
$$

Then $3'$ has extra entries of $\frac{2}{3}, -\frac{1}{3}, -\frac{1}{3}$, and $\bar{3}'$ has extra entries of $-\frac{2}{3}, \frac{1}{3}, \frac{1}{3}$. Thus, SU(3)'

6
(anti-)triplets of \( T_{1\pm}, T_{2\pm}, T_{40}, T_{4\pm}, T_{70}, \) and \( T_{7\pm} \) are

\[
T_{1\pm} : \ (1; 1, 3^{'})_{1/3} \rightarrow 1_1, \ 1_0, \ 1_0 \\
T_{2\pm} : \ (1; 1, 3^{'})_{1/3} \rightarrow 1_1, \ 1_0, \ 1_0 \\
T_{40} : \ 3 \cdot (1; 2; 1, 3^{'})_{1/6} \rightarrow 3 \cdot (1, 2)_{-1/2}, \ 3 \cdot (1, 2)_{1/2}, \ 3 \cdot (1, 2)_{1/2}; \\
\quad 3 \cdot (1; 1, 3^{'})_{-1/3} \rightarrow 3 \cdot 1_{-1}, \ 3 \cdot 1_0, \ 3 \cdot 1_0 \\
T_{4\pm} : \ 5 \cdot (1; 1, 3^{'})_{1/3} \rightarrow 5 \cdot 1_1, \ 5 \cdot 1_0, \ 5 \cdot 1_0 \\
\quad 2 \cdot (1; 1, 3^{'})_{-2/3} \rightarrow 2 \cdot 1_0, \ 2 \cdot 1_{-1}, \ 2 \cdot 1_{-1} \\
T_{70} : \ (1; 1, 3^{'})_{-2/3} \rightarrow 1_0, \ 1_{-1}, \ 1_{-1} \\
T_{7\pm} : \ (1; 1, 3^{'})_{-1/3} \rightarrow 1_{-1}, \ 1_0, \ 1_0
\]

(16)

Eq. (14) with (16) gives the SM quantum numbers. From these, we note that there is no exotics. Other exotics free orbifold compactifications \([6, 16]\) have \( E_8' \) sector contribution to \( Y \) as in the present case. But, we do not know whether this is a necessary condition for exotics free models or not.

### Table I

| \( P + [kV + ka] \) | \( \text{No.} \times \text{(Repts.)} \cdot Y_{Q_1 Q_2 Q_3 Q_4 Q_5} \) | \( \Gamma \) | Label |
|----------------------|-------------------------------------------------|--------|------|
| \( \frac{-1}{3} \frac{-1}{3} \frac{-1}{3} \frac{2}{3} \frac{2}{3} \frac{2}{3} \frac{-1}{3} \frac{-1}{3} 0 0 \) \( 0(8)^T_{T_{4\pm}} \) | \( 3 \cdot (3, 2)^L_{1/6} [0, 0, 0, 0] \) | 1 | \( q_1, q_2, q_3 \) |
| \( \frac{-1}{3} \frac{-1}{3} \frac{-1}{3} \frac{1}{3} \frac{1}{3} \frac{1}{3} \frac{-1}{3} \frac{-1}{3} \) \( 0(8)^T_{T_{3\pm}} \) | \( 2 \cdot (3, 1)^L_{-2/3} [-3, 3, 2, 0, 0] \) | 3 | \( u^c, e^c \) |
| \( \frac{1}{3} \frac{1}{3} \frac{1}{3} \frac{-1}{3} \frac{-1}{6} \frac{-1}{6} \frac{1}{3} \frac{1}{3} 0 0 \) \( 0(5)^T_{T_{10}} \) | \( (3, 1)^L_{1/3} [0, 6, -1, 5, 1] \) | 1 | \( t^c \) |
| \( \frac{1}{3} \frac{1}{3} \frac{1}{3} \frac{1}{3} \frac{1}{3} \frac{1}{3} \frac{-1}{3} \frac{-1}{3} \) \( 0(8)^T_{T_{10}} \) | \( (3, 1)^L_{1/3} [-3, 3, -2, 0, 0] \) | 1 | \( d^c, b^c \) |
| \( \frac{-1}{3} \frac{-1}{3} \frac{-1}{3} \frac{2}{3} \frac{2}{3} \frac{2}{3} \frac{-1}{3} \frac{-1}{3} 0 0 \) \( 0(8)^T_{T_{T_{10}}} \) | \( (1, 2)^L_{-1/2} [-6, 6, 0, 0, 0] \) | 1 | \( l_1, l_2, l_3 \) |
| \( 0 0 0 \frac{2}{3} \frac{2}{3} \frac{1}{3} \frac{1}{3} \frac{1}{3} \frac{1}{3} \frac{-1}{3} \frac{-1}{3} \) \( 0(8)^T_{T_{T_{10}}} \) | \( (1, 2)^L_{1/2} [0, 6, -1, 5, 1] \) | 0 | \( H_u \) |
| \( 0 0 0 \frac{2}{3} \frac{2}{3} \frac{1}{3} \frac{1}{3} \frac{1}{3} \frac{1}{3} \frac{-1}{3} \frac{-1}{3} \) \( 0(8)^T_{T_{T_{10}}} \) | \( (1, 2)^L_{-1/2} [-6, 0, -1, 5, 1] \) | -2 | \( H_d \) |
A. Three families with no exotics

Removing vectorlike representations and neutral singlets, we obtain the following chiral representations,

\[ T_{4, 7, 10} : 3 \cdot (3, 2)_{1/6}, \ 3 \cdot (3, 1)_{-2/3}, \ 3 \cdot (3, 1)_{1/3}, \ 3 \cdot (1, 2)_{-1/2}, \ 3 \cdot 1 \]  \tag{17} 

\[ T_{2, 7, 0} : 10', \ \overline{5}' \]  \tag{18} 

where \( 10' = (1; 10', 1) \) and \( \overline{5}' = (1; \overline{5}', 1) \). In Table I we list three families except the charged lepton singlets. Note that SU(3)_c triplets with underlined entries mean, for example,

\[ (\frac{-1}{3} \ -\frac{1}{3} \ -\frac{2}{3}) = (\frac{-1}{3} \ -\frac{1}{3} \ -\frac{2}{3}), (\frac{2}{3} \ -\frac{1}{3} \ \frac{1}{3}), (\frac{1}{6} \ -\frac{1}{6} \ \frac{1}{6}), (\frac{1}{6} \ -\frac{1}{6} \ \frac{1}{6}), (\frac{1}{6} \ -\frac{1}{6} \ \frac{1}{6}), (\frac{1}{6} \ -\frac{1}{6} \ \frac{1}{6}), (\frac{1}{6} \ -\frac{1}{6} \ \frac{1}{6}). \]

This is because of the asymmetrical simple roots of SU(3)_c in Eq. (5).

B. Matter parity

Let us define the U(1)_\Gamma charge as a linear combination of \( Q_1 - 5 \) of Eq. (11) and \( W_8 \). We choose its generator \( \Gamma \) such that the light quarks carry odd U(1)_\Gamma charges while Higgs doublets carry even U(1)_\Gamma charges. This is necessary to remove the baryon number violating \( u^c d^c d^c \) term. For the lepton number violation, the condition is not so strong and furthermore in our model there are so many possibilities in choosing the charged singlets \( e^c \), and here we do not discuss them. Then, one successful choice of \( \Gamma \) is

\[ \Gamma = \frac{1}{3}Q_2 + Q_3 + \overline{W}_8 \]  \tag{19} 

where

\[ \overline{W}_8 = (0^3 \ 1 \ 1 \ -2 \ 0^2)(0^8)' \]  

The \( \Gamma \) quantum numbers are also listed in Table I. Breaking U(1)_\Gamma by VEVs of even integer SM singlets, a discrete symmetry \( Z_2 \), which is called matter parity \( P \), survives,

\[ U(1)_\Gamma \rightarrow P. \]  \tag{20} 

Thus, looking at the light quarks only the dangerous term \( u^c d^c d^c \) is not allowed. However, we have to consider mixing of light quarks with heavy quarks which can be dangerous in principle [16]. In our model, there are ten quark flavors: six SM quarks and four extra \( Q_{em} = -\frac{1}{3} \) quarks denoted as \( 3 \cdot (D + \overline{D}) \) and \( (D' + \overline{D}') \). For quark mixing, we need to
consider $\overline{D}s$ and $\overline{D}'$ only. In Eq. (14), three $\overline{D}s$ (three out of five $(\overline{3}, 1)_{1/3}$) in $T_{4-}$ appear as $(\overline{3}, 1)_{1/3} [6, -6, 0, 0, 0]$ carrying $\Gamma = -2$ and $\overline{D}'$ in $T_{10}$ appears as $(\overline{3}, 1)_{1/3} [3, 3, 1, 5, 1]$ carrying $\Gamma = 2$. Therefore, if $P$ is not broken, light $\overline{d}c$ and heavy $\overline{D}s$ and $\overline{D}'$ can never mix and we achieve an exact matter parity $P$. But a successful matter parity assignment should not be in conflict with other phenomenological requirements. The most severe constraint comes from making exotic particles massive [16]. In passing, we point out that the other vectorlike particles, such as $D - \overline{D}, D' - \overline{D}'$, doublet pairs, and unit charge lepton pairs $E^- - \overline{E}^+$, are not so dangerous as exotics. Since our model does not include any exotics, we do not need VEVs of any odd $\Gamma$ singlets for obvious phenomenological reasons. A detailed study of singlet VEVs is outside of the scope of the present discussion, and will be presented elsewhere.

C. Higgs doublets

In Table II, we list all color singlet doublets, where we included lepton doublets in the last row. Higgs doublets form a vectorlike representation under the SM gauge group. So, they can be removed at the GUT scale in principle. One vectorlike pair $H_u + H_d$ is kept light for breaking the SU(2)×U(1)$_Y$ gauge symmetry at the electroweak scale. We choose the starred doublets to give large masses to $t$ and $b$ quarks. We choose $H_u$ such that the sum of the sector numbers in $q_3 t^c H_u$ adds up to $0 \mod 12$. Then, $H_u$ is chosen from $T_{10}$. For $b$ quark, a similar argument chooses one $(1, 2)_{-1/2}$ in $T_{4-}$ as $H_d$. These $H_u$ and $H_d$ are starred in Table II. However, note that this is just one illustration and another choice may well be possible depending on the Yukawa couplings and magnitudes of singlet VEVs.

D. The Lee-Weinberg model

This model is basically a string realization of the Lee-Weinberg model based on SU(3)$_c$×SU(3)$_W$×U(1) [17]. In the Lee-Weinberg model, one quark family consists of

$$3_{W,q} = \begin{pmatrix} d & u \\ D & \end{pmatrix}_L, \quad d_R, u_R, D_R$$

Thus, our model realizes just three left-handed quark triplets with no extra $3_W - \overline{3}_W$ quark pairs, and hence it is a minimal kind of Lee-Weinberg type models. Out of 21 left-handed $3_W$s and 21 left-handed $\overline{3}_W$s, 12 pairs form a vectorlike representations under the Lee-Weinberg
\[
\begin{array}{|c|c|}
\hline
P + n[V \pm a] & \Gamma \\
\hline
\left( \begin{array}{cccc}
\frac{1}{2} & \frac{1}{2} & \frac{1}{2} & \frac{1}{2} \\
0 & 0 & 0 & 0 \\
0 & 0 & 1 & 0 \\
0 & 0 & 1 & 0 \\
\end{array} \right)(0^8)'_{U_1} & -2 \\
\left( \begin{array}{cccc}
0 & 0 & 1 & 0 \\
0 & 0 & 1 & 0 \\
0 & 0 & 1 & 0 \\
0 & 0 & 0 & 0 \\
\end{array} \right)(0^8)'_{U_2} & 4 \\
\left( \begin{array}{cccc}
0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 \\
\end{array} \right)(0^8)'_{U_3} & 3 \\
\left( \begin{array}{cccc}
\frac{1}{2} & \frac{1}{2} & \frac{1}{2} & \frac{1}{2} \\
0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 \\
\end{array} \right)(0^8)'_{U_3} & 2 \\
\left( \begin{array}{cccc}
0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 \\
\end{array} \right)(0^8)'_{U_3} & -4 \\
\hline
\end{array}
\]

\[
\begin{array}{|c|c|}
\hline
\text{No. } \times (\text{Repts.}) & \text{Y}_{[Q_1,Q_2,Q_3,Q_4,Q_5]} \\
\hline
\left( \begin{array}{cccc}
1 & 2 & 1 & 2 \\
2 & L & 2 & L \\
2 & L & 2 & L \\
2 & L & 2 & L \\
\end{array} \right)_{\frac{1}{2}} & [9.3,-2;0,0] \\
\left( \begin{array}{cccc}
1 & 2 & 1 & 2 \\
2 & L & 2 & L \\
2 & L & 2 & L \\
2 & L & 2 & L \\
\end{array} \right)_{\frac{1}{2}} & [0.6,2;0,0] \\
\left( \begin{array}{cccc}
1 & 2 & 1 & 2 \\
2 & L & 2 & L \\
2 & L & 2 & L \\
2 & L & 2 & L \\
\end{array} \right)_{\frac{1}{2}} & [0.1,2;0,0] \\
\left( \begin{array}{cccc}
1 & 2 & 1 & 2 \\
2 & L & 2 & L \\
2 & L & 2 & L \\
2 & L & 2 & L \\
\end{array} \right)_{\frac{1}{2}} & [9.3,2;0,0] \\
\left( \begin{array}{cccc}
1 & 2 & 1 & 2 \\
2 & L & 2 & L \\
2 & L & 2 & L \\
2 & L & 2 & L \\
\end{array} \right)_{\frac{1}{2}} & [0,-6,-2;0,0] \\
\left( \begin{array}{cccc}
1 & 2 & 1 & 2 \\
2 & L & 2 & L \\
2 & L & 2 & L \\
2 & L & 2 & L \\
\end{array} \right)_{\frac{1}{2}} & [0,6,-1;5,1] \\
\left( \begin{array}{cccc}
1 & 2 & 1 & 2 \\
2 & L & 2 & L \\
2 & L & 2 & L \\
2 & L & 2 & L \\
\end{array} \right)_{\frac{1}{2}} & [0.2,0,-4] \\
\left( \begin{array}{cccc}
1 & 2 & 1 & 2 \\
2 & L & 2 & L \\
2 & L & 2 & L \\
2 & L & 2 & L \\
\end{array} \right)_{\frac{1}{2}} & [0,-6,1;5,3] \\
\left( \begin{array}{cccc}
1 & 2 & 1 & 2 \\
2 & L & 2 & L \\
2 & L & 2 & L \\
2 & L & 2 & L \\
\end{array} \right)_{\frac{1}{2}} & [0,-6,-1;5,3] \\
\left( \begin{array}{cccc}
1 & 2 & 1 & 2 \\
2 & L & 2 & L \\
2 & L & 2 & L \\
2 & L & 2 & L \\
\end{array} \right)_{\frac{1}{2}} & [0,0,0;0,0] \\
\left( \begin{array}{cccc}
1 & 2 & 1 & 2 \\
2 & L & 2 & L \\
2 & L & 2 & L \\
2 & L & 2 & L \\
\end{array} \right)_{\frac{1}{2}} & [0,0,0;0,0] \\
\left( \begin{array}{cccc}
1 & 2 & 1 & 2 \\
2 & L & 2 & L \\
2 & L & 2 & L \\
2 & L & 2 & L \\
\end{array} \right)_{\frac{1}{2}} & [0,0,0;0,0] \\
\left( \begin{array}{cccc}
1 & 2 & 1 & 2 \\
2 & L & 2 & L \\
2 & L & 2 & L \\
2 & L & 2 & L \\
\end{array} \right)_{\frac{1}{2}} & [0,0,0;0,0] \\
\hline
\end{array}
\]

TABLE II: Thirty-three color-singlet SU(2)\textsubscript{W} doublets which contain the leptons (the last row) and Higgs particles. The MSSM pair is starred.

gauge group.\textsuperscript{2} This is gleaned from the chiral representation \textsuperscript{17} that there remain three pairs of (3\textsubscript{c}, 3\textsubscript{W}). Thus, for SU(3)\textsubscript{W} anomaly cancellation, there must be nine 3\textsubscript{W}s, and the remaining 3\textsubscript{W} - 3\textsubscript{W} pairs must form a vectorlike representation. [We include the odd \Gamma Higgs pairs of Table III in the vectorlike representation.] Nine color-singlet 3\textsubscript{W}s contain three lepton doublets and three pairs of Higgs doublets. The electromagnetic charges of nine

\textsuperscript{2} The breaking scale of SU(3)\textsubscript{W} can be very low in principle, because the discrepancy in the numbers of multiplets between SU(3)\textsubscript{c} (ten flavors) and SU(3)\textsubscript{W} (twenty-one flavors) enables one to lower the breaking scale of SU(3)\textsubscript{W} while generating the difference of gauge couplings of SU(3)\textsubscript{c} and SU(3)\textsubscript{W}. But, we will not consider this possibility here.
\[ \bar{3}_W \text{s contain three } \bar{3}_{W,+} \text{ and six } \bar{3}_{W,0}, \text{ where} \]
\[ \bar{3}_{W,+} = \left( \begin{array}{c} \psi_1^+ \\ \psi^0 \\ \psi_2^+ \end{array} \right), \quad \bar{3}_{W,0} = \left( \begin{array}{c} \psi_1^0 \\ \psi^- \\ \psi_2^0 \end{array} \right)_L \]
\[ (22) \]

where \( \psi^{\text{sign}} \) denotes the integer electromagnetic charge of the field \( \psi \). In Eqs. (21) and (22), \( SU(2)_W \) doublets are pairs of \( u - d, \psi_2^+ - \psi_0, \) and \( \psi_2^0 - \psi^- \). Obviously, three lepton doublets of (17) must come from three \( \bar{3}_{W,0} \) s, and we are left with three pairs of \( \bar{3}_{W,0} - \bar{3}_{W,+} \).

E. The \( \mu \) term

A possible large \( \mu \) term arises from the coupling between three pairs of \( \bar{3}_{W,0} - \bar{3}_{W,+} \) as \( \epsilon_{\alpha\beta\gamma} \bar{3}_{W,0}^{\alpha} \bar{3}_{W,0}^{\beta} \bar{3}_{W,0}^{\gamma} \) where \( \alpha, \beta, \gamma \) are \( SU(3)_W \) indices. Suppose that \( SU(3)_W \) is broken by VEVs (typically of order \( V \)) of \( \psi_1^0 \) in \( \bar{3}_{W,0} \) (and also by \( \bar{3}_{W,0} \) in the removed vectorlike representation toward a \( D \)-flat condition). Then, the \( H_u - H_d \) type couplings arise from\(^3\)
\[ \epsilon_{\alpha\beta\gamma} \bar{3}_{W,I}^{\alpha} \bar{3}_{W,J}^{\beta} \bar{3}_{W,K}^{\gamma} \epsilon^{IJK} \sim V \epsilon_{\alpha\beta} \bar{3}_{W,I}^{\alpha} \bar{3}_{W,J}^{\beta} \epsilon^{IJ} \]
\[ (23) \]

where \( I, J, K \) are the Higgs family indices. For a general family coupling \( g^{IJK} \), due to \( \epsilon_{\alpha\beta\gamma} \) the symmetric part does not give an \( H_u - H_d \) coupling because it gives, \( \propto \bar{3}_{W,1} \bar{3}_{W,2} - \bar{3}_{W,2} \bar{3}_{W,1} = 0 \). Because of \( \epsilon^{IJ} \), the same Higgs family does not have the coupling and the \( 3 \times 3 \) \( H_u - H_d \) mass matrix is an antisymmetric one whose determinant is zero. Therefore, we obtain two massive Higgs doublet pairs and one massless Higgs doublet pair. Thus, there remains only one massless Higgs doublet pair, achieving the MSSM spectrum at low energy. In this scheme also, there are methods to generate an electroweak scale \( \mu \) term \([3, 4]\).

III. HIDDEN SECTOR \( SU(5)' \), GAUGE MEDIATION AND MESSENGERS

As shown in Table III there are \( SU(5)' \) fields. But some of these obtain masses by Yukawa couplings at the string scale. Below the string scale vectorlike pairs become massive by VEVs of singlets, and hence we consider only the chiral representations. We need the mass scale of the vectorlike pairs are much above the \( SU(5)' \) confining scale so that the SUSY breaking by \( 10' \) and \( 5' \) is intact.

\(^3\) Note that \( 3_W - \bar{3}_W \) coupling is not generating \( H_u - H_d \) terms since both \( H_u \) and \( H_d \) belong to \( 3_W \).
The quantum number of has been known that scale is achieved by VEVs of SM gauge singlet fields, breaking extra gauge symmetries. It′ and strength field W singlet combination in this uncalculable model can be parameterized by the gauge field •10 GMSB if the confining scale is below note that the singlet combination in Table III, we list all the SU(5)′ representations under SU(2)n×SU(5)′×SU(3)′. After removing vectorlike representations by = even integer singlets, the starred representations remain.

In Table III we list all the SU(5)′ non-singlet fields. From these, one can easily check that SU(5)′ gauge anomaly is absent. One conspicuous feature is that we obtained one 10′. Except 10′ of T2– and 5′ of T70, the rest 8 flavors form a vectorlike representation under SU(5)′×SU(2)n×U(1)Y. Removal of the eight flavors much above the SU(5)′ confining scale is achieved by VEVs of SM gauge singlet fields, breaking extra gauge symmetries. It has been known that 10′+5′ of a confining SU(5)′ breaks SUSY [11] and we achieve the GMSB if the confining scale is below 10^{12} GeV. Note that 10′0 and 5′0 do not carry any SU(3)c×SU(2)W×U(1)Y charge (which is emphasized by the subscript 0) and DSB by 10′0 and 5′0 does not break the SM gauge group.

Note that the singlet combination 10′10′10′5′ is not possible with one 10′. The SU(5)′ singlet combination in this uncalculable model can be parameterized by the gauge field strength field WαγWγ as discussed in [13]. The interaction between the messenger f and the hidden sector gauge fields can appear from string compactification as

\[ \mathcal{L} = \int d^2 \theta \left[ \xi(S_1, S_2, \cdots) f \bar{f} W^\alpha W^\gamma_\alpha + \eta(S_1, S_2, \cdots) f \bar{f} \right] + \text{h.c.} \]  

(24)

where we have in general the holomorphic functions \( \xi \) and \( \eta \) of singlet chiral fields, \( S_1, S_2, \cdots \). The quantum number of \( \xi(S_1, S_2, \cdots) f \bar{f} \) is the same as that of dilaton, where the H-
momentum of dilaton is $(0, 0, 0)$. On the other hand, the $H$-momenta of the superpotential term $\eta(S_1, S_2, \cdots) f \bar{T}$ should be $(-1, 1, 1)$. The $H$-momenta of the twisted sectors are given by [16, 18, 19]

\begin{align*}
U_1 & : (-1, 0, 0), \quad U_2 : (0, 1, 0), \quad U_3 : (0, 0, 1), \\
T_1 & : (\frac{-7}{12}, \frac{4}{12}, \frac{1}{12}), \quad T_2 : (\frac{-1}{6}, \frac{4}{6}, \frac{1}{6}), \quad T_3 : (\frac{-3}{4}, 0, \frac{1}{4}), \\
T_4 & : (\frac{-1}{3}, \frac{1}{3}, \frac{1}{3}, \frac{1}{3}), \quad \{T_5 : (\frac{1}{12}, \frac{-7}{12}, \frac{-7}{12}), \quad T_6 : (\frac{-1}{4}, 0, \frac{1}{4}), \\
T_7 & : (\frac{-1}{4}, \frac{1}{4}, \frac{7}{12}), \quad T_9 : (\frac{-1}{4}, 0, \frac{3}{4}),
\end{align*}

(25)

The Yukawa coupling $\eta(S_1, S_2, \cdots) f \bar{T}$ and the coefficient of $\mathcal{W}_c^\alpha \mathcal{W}_{\alpha}'$ must satisfy the modular invariance rule for the twisted sector fields $(z)$ multiplication,

$$\sum_z k(z) = 0 \mod 12, \quad \sum_z [km_f](z) = 0 \mod 3. \quad (26)$$

Consider, for example, the vectorlike colored particles appearing only in $T_{4-}$ with $Q_{em} = \mp \frac{1}{3}$: $f_3 = D, \bar{f}_3 = \bar{D}$, viz. Eq. [14]. We can consider the following gauge singlet combination multiplied to $\mathcal{W}_c^\alpha \mathcal{W}_{\alpha}'$, for example,

$$T_{4-} T_4 T_4 T_4 T_4 T_4 \sim \bar{f}_3 f_3 (T_4 T_7 T_4 T_4) \sim D_{-1/3} \bar{D}_{1/3}. \quad (27)$$

Similarly, $SU(2)_W$ doublet coupling $\mathcal{W}_c^\alpha \mathcal{W}_{\alpha}'$ can be considered. The product in (27), $T_{4-} T_4 T_4 T_4 T_4$, has the $H$-momentum $(-2, 2, 2)$, and hence we must multiply further singlets to make the sum of $H$-momenta be $(0, 0, 0)$. As shown in [16], usually we can achieve this, but here we do not elaborate the details. In this model, $f_3$ and $f_2$ denote the messenger through $SU(3)_c$ and the messenger through $SU(2)_W$, respectively. If needed, we can also consider $f_1$ (the messenger through $U(1)_Y$). Below, $f$ represents $f_3, f_2, \text{ or } f_1$.

From the above discussion, the fields of $f, \bar{f}$ and $\mathcal{W}_c^\alpha \mathcal{W}_{\alpha}'$ can have the following tree level Lagrangian,

$$\mathcal{L} = \int d^2 \theta \left[ \frac{1}{M^2} f \bar{f} \mathcal{W}_c^\alpha \mathcal{W}_{\alpha}' + M_f f \bar{T} \right] + \text{h.c.} \quad (28)$$

which is perturbative in origin. Here $M$ and $M_f$ are determined by the strength of coupling constant and VEVs of singlet fields appearing in $\xi$ and $\eta$ of Eq. [24]. Both of these parameters are assumed to be somewhat less than the string scale. The SUSY breaking through Eq. [28] has been discussed in [13] by introducing the messenger mass- and F-parameters

$$M_{\text{mess}} \approx M_f + \frac{A_3^3}{M^2}, \quad F_{\text{mess}} \approx \frac{A_4^4}{M^2}. \quad (29)$$
With this GMSB scenario, firstly the observable sector gaugino obtains mass of order
\[ \tilde{m}_{\text{SUSY}} \sim \frac{g^2}{16\pi^2} \frac{\Lambda_h^4}{M^2 M_{\text{mess}}} \]  
(30)
while the gravitino mass is around \( m_{3/2} \approx \Lambda_h^3/M_{\text{Pl}} \). To obtain 1 TeV gluino mass (but much smaller gravitino mass of order 0.2 GeV) with \( \alpha = \frac{1}{25} \) and \( \Lambda_h = 10^{12} \) GeV, for example, we need \( (M^2 M_{\text{mess}})^{1/3} \approx 1.5 \times 10^{14} \) GeV.

This leads us to consider the \( \mathcal{W}^\alpha \mathcal{W'}^\alpha \) couplings to \( H_u H_d \) and the observable sector Yukawa couplings \( W_Y \sim H_u q_i u_j^c + H_d q_i d_j^c \). Let us focus on the \( H_u H_d \) coupling. From the discussion with (23), the three pairs of Higgsinos form an antisymmetric mass matrix parametrized by \( A, B \) and \( C \) which are assumed to be large. The \( \int d^2 \theta H_u H_d \mathcal{W}^\alpha \mathcal{W'}^\alpha \) term would contribute to the Higgsino mass matrix and also to the soft \( B \) parameter matrix. The heavy pairs of \( H_u \) and \( H_d \) act as \( f_2 \) and \( \bar{f}_2 \). We are interested in the light \( H_u \) and \( H_d \) pair.

The Higgsino mass matrix and the \( B \) matrix take the following form,

\[ M_{\text{Higgsino}} = \begin{pmatrix}
0, & A + a, & B + b \\
-A - a, & 0, & C + c \\
-B - b, & -C - c, & 0
\end{pmatrix} \]
(31)
\[ B_{\text{soft}} = \mu \begin{pmatrix}
0, & a, & b \\
-a, & 0, & c \\
-b, & -c, & 0
\end{pmatrix} \]
(32)

where the parameters \( a, b, c \) in (31) get contribution from the hidden-sector gluino condensation while \( \mu(a, b, c) \) in (32) get contribution from the F-term of \( \mathcal{W}^\alpha \mathcal{W'}^\alpha \). If \( a : b : c = A : B : C \), then the light Higgsinos and light Higgs bosons are paired to constitute the Higgs multiplets of the MSSM. This proportionality can be achieved if the same singlet combination is multiplied to the six nonvanishing superpotential terms implied in (31) comprised of the \( H \)-momentum \((−1, 1, 1)\) to make the \( H \)-momentum \((0, 0, 0)\) of \( \xi f \bar{f} \) in (24). One may choose a vacuum so that such a condition is satisfied. The interaction \( \int d^2 \theta (\frac{1}{m} H_u q u^c + \cdots) \mathcal{W}^\alpha \mathcal{W'}^\alpha \) can be within a safe region of the gauge hierarchy solution. For example, the \( A \)-term estimated from this is \( A \approx \frac{\Lambda_h^4}{m} \) which can be of order \( 10^{-2} \) GeV – \( 10^6 \) GeV for \( \Lambda_h \sim 10^{10−12} \) GeV and \( m \sim 10^{14} \) GeV.

Finally, we comment on possible higher order terms in the Kähler potential. Even though all the important hidden sector matter \( 10' \) does not appear in the superpotential, it can
appear in the Kähler potential. Possible terms of the form $10'10'^*f f^*/M_K^2$ might appear. The higher order Kähler terms was calculated for the compactification $T^6 = (T_2)^3$ (with the volume moduli $T$s and the complex structure moduli $C$s) in Ref. [20] for two matter fields $Q_\alpha$,

$$K_{\text{matter}} = \prod_{i=1}^{3} (T_i + \overline{T}_i)^{n^i_\alpha} \prod_{m=1}^{h_{(2,1)}} (C_m + \overline{C}_m)^{l_\alpha} |Q_\alpha|^2$$

where $n^i_\alpha$ and $h_{2,1} = 1$ (for our $\mathbb{Z}_{12-I}$) are the modular weight and a Hodge number, respectively. Also, $l_\alpha$ is an integer. The term $10'10'^*f f^*/M_K^2$ is not appearing in the above expression, and at present there does not exist a $K_{\text{matter}}$ calculation for four matter fields of our interest. Even if it appears, the mass suppression scale $M_K$ is expected to be of order the string scale and hence is much larger than $M$ appearing in Eq. (28) toward the GMSB scenario. However, if it appears with the same order of the suppression factor as in Eq. (28), the idea of our GMSB is not successful phenomenologically. We may need $M^2/M_K^2 < 0.03$.

IV. CONCLUSION

We have shown that there exists a possibility of the hidden sector SU(5)$'$ with $10'_0$ plus $\overline{5}'_0$ matter below the GUT scale so that a GMSB at the stable vacuum is successful. Toward achieving the needed coupling constant $\alpha'_5$ of the hidden sector at the GUT scale, we may need different compactification radii for the three tori [6]. The model is very interesting in that it contains three MSSM families without any exotics. We find a desirable U(1)$_I$ gauge symmetry whose $\mathbb{Z}_2$ discrete group can be a matter parity $P$ or $R$-parity. Due to our Lee-Weinberg type model, there remains only one light pair of Higgs doublets, achieving the MSSM spectrum. On the other hand, the weak mixing angle at the unification scale is not $\frac{\pi}{8}$. Various mass scales in addition to the different compactification radii may enable us to fit the mixing angle to the observed one at the electroweak scale. A detail analysis of the model for the $R$-parity problem, weak mixing angle, compactification radii, $D$ and $F$ flat directions, and Yukawa couplings will be discussed elsewhere.
Acknowledgments

I thank K.-S. Choi, I. W. Kim and B. Kyae for helpful discussions. This work is supported in part by the KRF Grants, No. R14-2003-012-01001-0 and No. KRF-2005-084-C00001.

[1] F. Gabbiani, E. Gabrielli, A. Masiero, and L. Silvestrini, Nucl. Phys. B477 (1996) 321.
[2] For a review, see, D. J. H. Chung, L. L. Everett, G. L. Kane, S. F. King, J. D. Lykken, and L. T. Wang, Phys. Rep. 407 (2005) 1.
[3] J. E. Kim and H. P. Nilles, Phys. Lett. B138 (1984) 150.
[4] G. Giudice and A. Masiero, Phys. Lett. B206 (1988) 480; E. J. Chun, J. E. Kim and H. P. Nilles, Nucl. Phys. B370 (1992) 105.
[5] M. Dine and A. E. Nelson, Phys. Rev. D48 (1993) 1277; M. Dine, A. E. Nelson and Y. Shirman, Phys. Rev. D51 (1995) 1362. For reviews, see, K. A. Intrilligator and N. Seiberg, Nucl. Phys. Proc. Suppl. 45BC (1996) 1; Y. Shadmi and Y. Shirman, Rev. Mod. Phys. 72 (2000) 25.
[6] J. E. Kim, Phys. Lett. B651 (2007) 407 [arXiv:0706.0293].
[7] K. A. Intrilligator, N. Seiberg and D. Shih, JHEP 0604 (2006) 021 [hep-th/0602239].
[8] H. Murayama and Y. Nomura, Phys. Rev. Lett. 98 (2007) 151803 [hep-ph/0612186].
[9] R. Kitano, H. Ooguri, and Y. Ookouchi, Phys. Rev. D75 (2007) 045022 [hep-th/0612139]; R. Argurio, M. Bertolini, S. Franco, and S. Kachru, [hep-th/0703236]; A. Amariti, L. Girardello, and A. Mariotti, [hep-th/0701121]; M. Dine and J. Mason, [hep-ph/0611312]; C. Csaki, Y. Shirman and J. Terning, [hep-ph/0612241]; O. Aharony and N. Seiberg, JHEP 0702 (2007) 054 [hep-ph/0612308]; I. Affleck, M. Dine and N. Seiberg, Nucl. Phys. B256 (1985) 557; H. Murayama, Phys. Lett. B355 (1995) 187.
[10] I. Affleck, M. Dine and N. Seiberg, Nucl. Phys. B256 (1985) 557; H. Murayama, Phys. Lett. B355 (1995) 187.
[11] G. Veneziano, Phys. Lett. B128 (1984) 199; I. Affleck, M. Dine and N. Seiberg, Nucl. Phys. B256 (1985) 557; E. Poppitz and S. P. Trivedi, Phys. Lett. B401 (1997) 38.
[12] J. E. Kim and B. Kyae, Nucl. Phys. B770 (2007) 47 [hep-th/0608086].
[13] H. Murayama and Y. Nomura, Phys. Rev. D75 (2007) 095011 [hep-ph/0701231].
[14] D. E. Diaconescu, B. Florea, S. Kachru, and P. Svrcek, JHEP \textbf{0602} (2006) 020 [hep-th/0512170].

[15] L. J. Dixon, J. A. Harvey, C. Vafa and E. Witten, Nucl. Phys. \textbf{B261} (1985) 678; Nucl. Phys. \textbf{B274} (1986) 285; L. Ibanez, H. P. Nilles and F. Quevedo, Phys. Lett. \textbf{B187} (1987) 25.

[16] J. E. Kim, J.-H. Kim and B. Kyae, JHEP \textbf{0706} (2007) 034 [hep-ph/0702278].

[17] B. W. Lee and S. Weinberg, Phys. Rev. Lett. \textbf{38} (1977) 1237; B. W. Lee and R. E Shrock, Phys. Rev. \textbf{D17} (1977) 2410.

[18] K.-S. Choi and J. E. Kim, \textit{Quarks and Leptons from Orbifolded Superstring} (Springer-Verlag, Heidelberg, Germany, 2006).

[19] Y. Katsuki, Y. Kawamura, T. Kobayashi, N. Otsubo, Y. Ono, and K. Tanioka, Nucl. Phys. \textbf{B341} (1990) 611.

[20] L. E. Ibanez and D. Lüst, Nucl. Phys. \textbf{B382} (1992) 305; D. Bailin and A. Love, Phys. Lett. \textbf{B288} (1992) 263.

[21] M. Dine, Y. Nir and Y. Shirman, Phys. Rev. \textbf{D55} (1997) 1501.