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

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Abstract

We show that it is possible to introduce the confining hidden sector gauge group $SU(5)'$ with the chiral matter $\mathbf{10}'_0$ plus $\bar{\mathbf{5}}'_0$, which are neutral under the standard model gauge group, toward a gauge mediated supersymmetry breaking (GMSB) in a $\mathbf{Z}_{12-I}$ orbifold compactification of $\mathrm{E}_8 \times \mathrm{E}'_8$ 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 μ problem.

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1. 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 μ problem [3], etc. The hierarchical magnitude is worst in the μ 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) [7–9], 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 $N_c + 1 \leqslant N_f < \frac{3}{2}N_c$. This mechanism is easily applicable to $SU(5)'$ models with 6 or 7 flavors, which can be realized in string compactifications [6]. 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 Letter, therefore, we look for a GMSB spectrum in the orbifold compactification of the $\mathrm{E}_8 \times \mathrm{E}'_8$ heterotic string with three families, trying to satisfy all obvious phenomenological requirements.

The well-known DSB models are an $SO(10)'$ model with $\mathbf{16}'$ or $\mathbf{16}' + \mathbf{10}'$ [10], and an $SU(5)'$ model with $\mathbf{10}' + \bar{\mathbf{5}}'$ [11]. It is known that GMSB with $\mathbf{16}' + \mathbf{10}'$ 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 [6]. To solve

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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]

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

where Q is a hidden sector quark and f is a messenger. It is possible because the vectorlike representations, for example six or seven $(Q+\bar{Q})$, are present and the $Q\bar{Q}f\bar{f}$ 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 $\mathbf{10}'+\bar{\mathbf{5}}'$ of $SU(5)'$ does not have such a simple singlet direction in terms of chiral fields. For example, the term $\epsilon_{ijklm}10^{ij}10^{kl}10^{mn}\bar{5}_n = 0$ since taking $n=1$ without generality it is proportional to $\epsilon_{1jklm}10^{1j}10^{kl}10^{m1}\bar{5}_1$ which can be shown to be vanishing using the antisymmetric symbol ϵ. The singlet combination is possible in terms of the chiral gauge field strength, $\mathcal{W}'^{\alpha}\mathcal{W}'_{\alpha}$. It is pointed out that the F-term of this singlet combination can trigger the SUSY breaking to low energy [13],

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

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 [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 $\mathbf{Z}_{12-I}$ 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 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 μ 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.

2. A $\mathbf{Z}_{12-I}$ model

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

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

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 $\mathbf{Z}_{12-I}$ in the $\mathrm{E}_8\times\mathrm{E}_8'$ space in terms of the shift vector V and the Wilson line a_3 as

$$V = \frac{1}{12}(6\,6\,6\,2\,2\,2\,3\,3)(3\,3\,3\,3\,3\,1\,1\,1)', \tag{2}$$

$$a_3 = \frac{1}{3}(1\,1\,2\,0\,0\,0\,0\,0)(0\,0\,0\,0\,0\,1\,1\,-2)'. \tag{3}$$

(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 integer,

$$\begin{aligned} &SU(3)_c\times SU(3)_W\times SU(2)_n\times U(1)_a\times U(1)_b\times U(1)_c \\ &\quad\times\left[SU(5)'\times SU(3)'\times U(1)'^2\right]. \end{aligned} \tag{4}$$

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

$$SU(3)_c\colon \begin{cases} \alpha_1 = (1\,-1\,0\,0\,0\,0\,0\,0), \\ \alpha_2 = (0\,1\,1\,0\,0\,0\,0\,0), \end{cases} \tag{5}$$

$$SU(2)_W\colon \{\alpha_1 = (0\,0\,0\,1\,-1\,0\,0\,0), \tag{6}$$

$$SU(2)_n\colon \alpha_1 = (0\,0\,0\,0\,0\,0\,1\,-1). \tag{7}$$

The hypercharge direction is the combination of $U(1)$s of Eq. (4) and some generators of non-Abelian groups

$$\begin{aligned} Y &= Y_{\text{Abel}} + \frac{1}{\sqrt{3}}W_8 + F_3 - \frac{1}{\sqrt{3}}F_8 \\ &= \tilde{Y} + F_3 - \frac{1}{\sqrt{3}}F_8, \end{aligned} \tag{8}$$

where

$$Y_{\text{Abel}} = Y_8 + Y_8', \tag{9}$$

and W_8, F_3, F_8 are non-Abelian 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 $SU(3)_W$, $(\frac{1}{2\sqrt{3}}, \frac{1}{2\sqrt{3}}, -\frac{1}{\sqrt{3}})$, and F_3 and F_8 are the third and the eighth generators of $SU(3)'$, $(\frac{1}{2}, -\frac{1}{2}, 0)$ and $(\frac{1}{2\sqrt{3}}, \frac{1}{2\sqrt{3}}, -\frac{1}{\sqrt{3}})$, respectively. We find that exotics cannot be made vectorlike if we do not include Y'. $\tilde{Y}$ is defined as

$$\begin{aligned} \tilde{Y} &= Y_{\text{Abel}} + \frac{1}{\sqrt{3}}W_8 \\ &= \left(\frac{1}{6}\,\frac{1}{6}\,\frac{-1}{6}\,0\,0\,\frac{-1}{2}\,\frac{-1}{2}\,\frac{-1}{2}\right)\left(\frac{1}{2}\,\frac{1}{2}\,\frac{1}{2}\,\frac{1}{2}\,\frac{1}{2}\,\frac{-1}{6}\,\frac{-1}{6}\,\frac{-1}{6}\right)'. \end{aligned} \tag{10}$$

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

$$Q_1 = (6\,6\,-6\,0\,0\,0\,0\,0)(0\,0\,0\,0\,0\,0\,0\,0)',$$

$$Q_2 = (0\,0\,0\,6\,6\,6\,0\,0)(0\,0\,0\,0\,0\,0\,0\,0)',$$

[1] This form has been considered by many [9], in particular in [8].

$$Q_3 = (0\,0\,0\,0\,0\,0\,2\,2)(0\,0\,0\,0\,0\,0\,0\,0)',$$
$$Q_4 = (0\,0\,0\,0\,0\,0\,0\,0)(4\,4\,4\,4\,4\,0\,0\,0)',$$
$$Q_5 = (0\,0\,0\,0\,0\,0\,0\,0)(0\,0\,0\,0\,0\,4\,4\,4)'. \quad (11)$$

(b) *Matter representations*: Now there is a standard method to obtain the massless spectrum in $\mathbf{Z}_{12-I}$ orbifold models. The spectra in the untwisted sectors U_1, U_2, and U_3, and twisted sectors, $T1_{0,+,-}, T2_{0,+,-}, T3, T4_{0,+,-}, T5_{0,+,-}$, and $T6$, are easily obtained [16]. The representations are denoted as

$$\left[\mathbf{SU}(3)_c, \mathbf{SU}(2)_W; \mathbf{SU}(5)', SU(3)'\right]_{\tilde{Y}}, \quad (12)$$

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

$$\left(\mathbf{SU}(3)_c, \mathbf{SU}(2)_W\right)_Y. \quad (13)$$

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

$$\begin{aligned}
&U_1\!: (\mathbf{1},\mathbf{2})_{1/2}, 2\cdot(\mathbf{1},\mathbf{2})_{-1/2}, \mathbf{1}_1, 2\cdot\mathbf{1}_0,\\
&U_2\!: (\mathbf{1},\mathbf{2})_{-1/2}, \mathbf{1}_0,\\
&U_3\!: (\mathbf{1},\mathbf{2})_{-1/2}, 2\cdot(\mathbf{1},\mathbf{2})_{1/2}, 2\cdot\mathbf{1}_1,\\
&T_{1_0}\!: (\bar{\mathbf{3}},\mathbf{1})_{1/3},\ (\mathbf{1},\mathbf{2})_{1/2},\ 3\cdot\mathbf{1}_1,\ 2\cdot\mathbf{1}_0,\\
&T_{1_-}\!: (\mathbf{1};\bar{\mathbf{5}}',\mathbf{1})_0, (\mathbf{1};\mathbf{1},\mathbf{3}')_{1/3}, 2\cdot\mathbf{1}_{-1},\\
&T_{2_0}\!: (\bar{\mathbf{3}},\mathbf{1})_{1/3}, (\mathbf{1},\mathbf{2})_{-1/2}, 3\cdot\mathbf{1}_0,\\
&T_{2_+}\!: (\mathbf{1};\mathbf{10}',\mathbf{1})_0, (\mathbf{1};\mathbf{1},\mathbf{3}')_{1/3}, 4\cdot\mathbf{1}_0,\\
&T_3\!: 2\cdot(\mathbf{1};\mathbf{5}',\mathbf{1})_0, 2\cdot(\mathbf{1};\bar{\mathbf{5}}',\mathbf{1})_0,\\
&\quad (2_L+1_R)(\mathbf{1},\mathbf{2})_{1/2}, (1_L+2_R)(\mathbf{1},\mathbf{2})_{1/2},\\
&\quad (2_L+1_R)\mathbf{1}_1, 3\cdot\mathbf{1}_0, (6_L+6_R)\cdot\mathbf{1}_1,\\
&T_{4_0}\!: 3\cdot(\mathbf{1},\mathbf{2};\mathbf{1},\bar{\mathbf{3}}')_{1/6}, 3\cdot(\mathbf{1};\mathbf{1},\bar{\mathbf{3}}')_{-1/3},\\
&T_{4_+}\!: 5\cdot(\mathbf{1};\mathbf{1},\mathbf{3}')_{1/3}, 2\cdot(\mathbf{1};\mathbf{1},\mathbf{3}')_{-2/3},\\
&T_{4_-}\!: 3\cdot(\mathbf{3},\mathbf{2})_{1/6}, 2\cdot(\bar{\mathbf{3}},\mathbf{1})_{-2/3}, 5\cdot(\bar{\mathbf{3}},\mathbf{1})_{1/3}, 3\cdot(\mathbf{3},\mathbf{1})_{-1/3},\\
&\quad 5\cdot(\mathbf{1},\mathbf{2})_{-1/2}, 2\cdot(\mathbf{1},\mathbf{2})_{1/2}, 2\cdot\mathbf{1}_1, 12\cdot\mathbf{1}_0, 12\cdot\mathbf{1}_0,\\
&T_{7_0}\!: (\mathbf{1};\bar{\mathbf{5}}',\mathbf{1})_0, (\mathbf{1};\mathbf{1},\mathbf{3}')_{-2/3},\\
&T_{7_+}\!: (\bar{\mathbf{3}},\mathbf{1})_{-2/3}, (\mathbf{3},\mathbf{1})_{-1/3}, 2\cdot(\mathbf{1},\mathbf{2})_{-1/2}, \mathbf{1}_0, 3\cdot\mathbf{1}_{-1},\\
&T_{7_-}\!: (\mathbf{1};\mathbf{5}',\mathbf{1})_0, (\mathbf{1};\mathbf{1},\bar{\mathbf{3}}')_{-1/3}, 2\cdot\mathbf{1}_1,\\
&T_6\!: 3\cdot(\mathbf{1};\mathbf{5}',\mathbf{1})_0, 3\cdot(\mathbf{1};\bar{\mathbf{5}}',\mathbf{1})_0, 2\cdot(\mathbf{1};\mathbf{5}',\mathbf{1})_1,\\
&\quad 2\cdot(\mathbf{1};\bar{\mathbf{5}}',\mathbf{1})_{-1},
\end{aligned} \quad (14)$$

where $\mathbf{1} = (\mathbf{1},\mathbf{1},\mathbf{1};\mathbf{1},\mathbf{1})$. Breaking $SU(3)'$, we assign

$$F_3 = \left(\frac{1}{2}, -\frac{1}{2}, 0\right), \qquad \frac{1}{\sqrt{3}} F_8 = \left(\frac{1}{6}, \frac{1}{6}, -\frac{1}{3}\right). \quad (15)$$

Then $\mathbf{3}'$ has extra entries of $\frac{2}{3}, -\frac{1}{3}, -\frac{1}{3}$, and $\bar{\mathbf{3}}'$ has extra entries of $-\frac{2}{3}, \frac{1}{3}, \frac{1}{3}$. Thus, $SU(3)'$ (anti-)triplets of $T_{1_-}, T_{2_+}, T_{4_0}, T_{4_+}, T_{7_0}$ and T_{7_-} are

$$\begin{aligned}
&T_{1_-}\!: (\mathbf{1};\mathbf{1},\mathbf{3}')_{1/3} \to \mathbf{1}_1,\ \mathbf{1}_0,\ \mathbf{1}_0,\\
&T_{2_+}\!: (\mathbf{1};\mathbf{1},\mathbf{3}')_{1/3} \to \mathbf{1}_1,\ \mathbf{1}_0,\ \mathbf{1}_0,\\
&T_{4_0}\!: 3\cdot(\mathbf{1},\mathbf{2};\mathbf{1},\bar{\mathbf{3}}')_{1/6} \to 3\cdot(\mathbf{1},\mathbf{2})_{-1/2},\ 3\cdot(\mathbf{1},\mathbf{2})_{1/2},\\
&\qquad 3\cdot(\mathbf{1},\mathbf{2})_{1/2},\\
&\quad 3\cdot(\mathbf{1};\mathbf{1},\bar{\mathbf{3}}')_{-1/3} \to 3\cdot\mathbf{1}_{-1},\ 3\cdot\mathbf{1}_0,\ 3\cdot\mathbf{1}_0,\\
&T_{4_+}\!: 5\cdot(\mathbf{1};\mathbf{1},\mathbf{3}')_{1/3} \to 5\cdot\mathbf{1}_1,\ 5\cdot\mathbf{1}_0,\ 5\cdot\mathbf{1}_0,\\
&\quad 2\cdot(\mathbf{1};\mathbf{1},\mathbf{3}')_{-2/3} \to 2\cdot\mathbf{1}_0,\ 2\cdot\mathbf{1}_{-1},\ 2\cdot\mathbf{1}_{-1},\\
&T_{7_0}\!: (\mathbf{1};\mathbf{1},\mathbf{3}')_{-2/3} \to \mathbf{1}_0,\ \mathbf{1}_{-1},\ \mathbf{1}_{-1},\\
&T_{7_-}\!: (\mathbf{1};\mathbf{1},\bar{\mathbf{3}}')_{-1/3} \to \mathbf{1}_{-1},\ \mathbf{1}_0,\ \mathbf{1}_0.
\end{aligned} \quad (16)$$

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.

2.1. Three families with no exotics

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

$$T_{4_-,7_+,1_0}\!: 3\cdot(\mathbf{3},\mathbf{2})_{1/6}, 3\cdot(\mathbf{3},\mathbf{1})_{-2/3},$$
$$\quad 3\cdot(\mathbf{3},\mathbf{1})_{1/3}, 3\cdot(\mathbf{1},\mathbf{2})_{-1/2}, 3\cdot\mathbf{1}_1, \quad (17)$$

$$T_{2_+,7_0}\!: \mathbf{10}_0', \bar{\mathbf{5}}_0', \quad (18)$$

where $\mathbf{10}_0' = (\mathbf{1};\mathbf{10}',\mathbf{1})_0$ and $\bar{\mathbf{5}}_0' = (\mathbf{1};\bar{\mathbf{5}}',\mathbf{1})_0$. In Table 1, we list three families except the charged lepton singlets. Note that $SU(3)_c$ triplets with underlined entries mean, for example, $(\underline{\frac{-1}{3}\,\frac{-1}{3}\,\frac{-2}{3}}) = (\frac{-1}{3}\,\frac{-1}{3}\,\frac{-2}{3}), (\frac{2}{3}\,\frac{-1}{3}\,\frac{1}{3}), (\frac{-1}{3}\,\frac{2}{3}\,\frac{1}{3})$, and $(\underline{\frac{1}{6}\,\frac{1}{6}\,\frac{5}{6}}) = (\frac{1}{6}\,\frac{1}{6}\,\frac{5}{6}), (\frac{1}{6}\,\frac{-5}{6}\,\frac{-1}{6}), (\frac{-5}{6}\,\frac{1}{6}\,\frac{-1}{6})$. This is because of the asymmetrical simple roots of $SU(3)_c$ in Eq. (5).

2.2. 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 Γ 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 Γ is

$$\Gamma = \frac{1}{3} Q_2 + Q_3 + \tilde{W}_8, \quad (19)$$

where

$$\tilde{W}_8 = \left(0^3 1 1 -2\, 0^2\right)\left(0^8\right)'.$$

The Γ quantum numbers are also listed in Table 1. Breaking $U(1)_\Gamma$ by VEVs of even integer SM singlets, a discrete symmetry $\mathbf{Z}_2$, which is called *matter parity* P, survives,

$$U(1)_\Gamma \to P. \quad (20)$$

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 $Q_{\rm em} = -\frac{1}{3}$ quarks denoted as $3\cdot(D+\bar{D})$

Table 1
Three families of quarks and leptons and a pair of Higgs doublets. We do not list singlet leptons since there are many possibilities

$P+[kV+ka]$	No.$\times$(Repts.)$_{Y[Q_1,Q_2,Q_3,Q_4,Q_5]}$	Γ	Label
$(\underline{\frac{-1}{3}\frac{-1}{3}\frac{-2}{3}}\,\underline{\frac{2}{3}\frac{-1}{3}}\frac{-1}{3}00)(0^8)'_{T_{4_-}}$	$3\cdot(\mathbf{3},\mathbf{2})^L_{1/6[0,0,0;0,0]}$	1	q_1,q_2,q_3
$(\underline{\frac{1}{6}\frac{1}{6}\frac{5}{6}}\frac{1}{6}\frac{1}{6}\frac{1}{6}\frac{1}{2}\frac{1}{2})(0^8)'_{T_{4_-}}$	$2\cdot(\bar{\mathbf{3}},\mathbf{1})^L_{-2/3[-3,3,2;0,0]}$	3	u^c,c^c
$(\underline{\frac{-1}{3}\frac{-1}{3}\frac{-2}{3}}\frac{1}{3}\frac{1}{3}\frac{1}{3}\frac{-1}{4}\frac{-1}{4})(\frac{1}{4}^5\frac{1}{12}\frac{1}{12}\frac{1}{12})'_{T_{7_+}}$	$(\bar{\mathbf{3}},\mathbf{1})^L_{-2/3[0,6,-1;5,1]}$	1	t^c
$(\underline{\frac{1}{2}\frac{1}{2}\frac{1}{2}}\frac{-1}{6}\frac{-1}{6}\frac{-1}{6}00)(0^5\frac{-1}{3}\frac{-1}{3}\frac{-1}{3})'_{T_{2_0}}$	$(\bar{\mathbf{3}},\mathbf{1})^L_{1/3[3,-3,0;0,-4]}$	-1	d^c
$(\underline{\frac{1}{6}\frac{1}{6}\frac{5}{6}}\frac{1}{6}\frac{1}{6}\frac{1}{6}\frac{-1}{2}\frac{-1}{2})(0^8)'_{T_{4_-}}$	$2\cdot(\bar{\mathbf{3}},\mathbf{1})^L_{1/3[-3,3,-2;0,0]}$	1	s^c,b^c
$(\frac{-1}{3}\frac{-1}{3}\frac{1}{3}\underline{\frac{2}{3}\frac{-1}{3}}\frac{2}{3}00)(0^8)'_{T_{4_-}}$	$(\mathbf{1},\mathbf{2})^L_{-1/2[-6,6,0;0,0]}$	1	l_1,l_2,l_3
$(000\underline{\frac{2}{3}\frac{-1}{3}}\frac{2}{3}\frac{-1}{4}\frac{-1}{4})(\frac{1}{4}^5\frac{1}{12}\frac{1}{12}\frac{1}{12})'_{T_{1_0}}$	$(\mathbf{1},\mathbf{2})^L_{1/2[0,6,-1;5,1]}$	0	H_u
$(\frac{-1}{3}\frac{-1}{3}\frac{1}{3}\underline{\frac{1}{3}\frac{-2}{3}}\frac{1}{3}\frac{-1}{4}\frac{-1}{4})(\frac{1}{4}^5\frac{1}{12}\frac{1}{12}\frac{1}{12})'_{T_{7_+}}$	$(\mathbf{1},\mathbf{2})^L_{-1/2[-6,0,-1;5,1]}$	-2	H_d

and $(D'+\bar{D}')$. For quark mixing, we need to consider $\bar{D}$s and $\bar{D}'$ only. In Eq. (14), three $\bar{D}$s (three out of five $(\bar{\mathbf{3}},\mathbf{1})_{1/3}$s) in T_{4_-} appear as $(\bar{\mathbf{3}},\mathbf{1})_{1/3[6,-6,0;0,0]}$ carrying $\Gamma=-2$ and $\bar{D}'$ in T_{1_0} appears as $(\bar{\mathbf{3}},\mathbf{1})_{1/3[3,3,1;5,1]}$ carrying $\Gamma=2$. Therefore, if P is not broken, light d^c and heavy $\bar{D}$s and $\bar{D}'$ 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 $D-\bar{D}$, $D'-\bar{D}'$, doublet pairs, and unit charge lepton pairs $E^- - \bar{E}^+$, are not so dangerous as exotics. Since our model does not include any exotics, we do not need VEVs of any odd Γ 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.

2.3. Higgs doublets

In Table 2, 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 $SU(2)\times U(1)_Y$ 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 $q_3t^cH_u$ adds up to 0 mod 12. Then, H_u is chosen from T_{1_0}. For b quark, a similar argument chooses one $(\mathbf{1},\mathbf{2})_{-1/2}$ in T_{4_-} as H_d. These H_u and H_d are starred in Table 2. 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.

2.4. The Lee–Weinberg model

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

$$\mathbf{3}_{W,q}=\begin{pmatrix} d & & u \\ & D & \end{pmatrix}_L,\qquad d_R,\ u_R,\ D_R. \tag{21}$$

Thus, our model realizes just three left-handed quark triplets with *no extra* $\mathbf{3}_W-\bar{\mathbf{3}}_W$ *quark pairs*, and hence it is a minimal kind of Lee–Weinberg type models. Out of 21 left-handed $\mathbf{3}_W$s and 21 left-handed $\bar{\mathbf{3}}_W$s, 12 pairs form a vectorlike representations under the Lee-Weinberg gauge group.[2] This is gleaned from the chiral representation (17) that there remain three pairs of $(\mathbf{3}_c,\mathbf{3}_W)$. Thus, for $SU(3)_W$ anomaly cancellation, there must be nine $\bar{\mathbf{3}}_W$s, and the remaining $\mathbf{3}_W-\bar{\mathbf{3}}_W$ pairs must form a vectorlike representation. [We include the odd Γ Higgs pairs of Table 2 in the vectorlike representation.] Nine color-singlet $\bar{\mathbf{3}}_W$s contain three lepton doublets and three pairs of Higgs doublets. The electromagnetic charges of nine $\bar{\mathbf{3}}_W$s contain three $\bar{\mathbf{3}}_{W,+}$ and six $\bar{\mathbf{3}}_{W,0}$, where

$$\bar{\mathbf{3}}_{W,+}=\begin{pmatrix} & \psi_1^+ & \\ \psi^0 & & \psi_2^+ \end{pmatrix}_L,\qquad \bar{\mathbf{3}}_{W,0}=\begin{pmatrix} & \psi_1^0 & \\ \psi^- & & \psi_2^0 \end{pmatrix}_L, \tag{22}$$

where ψ^{sign} denotes the integer electromagnetic charge of the field ψ. 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 $\bar{\mathbf{3}}_{W,0}$s, and we are left with three pairs of $\bar{\mathbf{3}}_{W,0}-\bar{\mathbf{3}}_{W,+}$.

2.5. The μ term

A possible large μ term arises from the coupling between three pairs of $\bar{\mathbf{3}}_{W,0}-\bar{\mathbf{3}}_{W,+}$ as $\epsilon_{\alpha\beta\gamma}\bar{\mathbf{3}}_W^\alpha\bar{\mathbf{3}}_W^\beta\bar{\mathbf{3}}_W^\gamma$ where α,β,γ are $SU(3)_W$ indices. Suppose that $SU(3)_W$ is broken by VEVs (typically of order V) of ψ_1^0 in $\bar{\mathbf{3}}_{W,0}$ (and also by $\mathbf{3}_{W,0}$ in the removed vectorlike representation toward a D-flat condition).

[2] The breaking scale of $SU(3)_W$ can be very low in principle, because the discrepancy in the numbers of multiplets between $SU(3)_c$ (ten flavors) and $SU(3)_W$ (twenty-one flavors) enables one to lower the breaking scale of $SU(3)_W$ while generating the difference of gauge couplings of $SU(3)_c$ and $SU(3)_W$. But, we will not consider this possibility here.

Table 2
Thirty-three color-singlet $SU(2)_W$ doublets which contain the leptons (the last row) and Higgs particles. The MSSM pair is starred

$P+n[V\pm a]$	Γ	No.$\times$(Repts.)$_{Y[Q_1,Q_2,Q_3,Q_4,Q_5]}$
$(\frac{1}{2}\,\frac{1}{2}\,\frac{-1}{2}\,\underline{\frac{1}{2}\,\frac{-1}{2}}\,\frac{1}{2}\,\frac{-1}{2}\,\frac{-1}{2})(0^8)'_{U_1}$	-2	$(\mathbf{1},\mathbf{2})^L_{\frac{1}{2}[9,3,-2;0,0]}$
$(000\underline{10}0\underline{10})(0^8)'_{U_1}$	4	$(\mathbf{1},\mathbf{2},\mathbf{2})^L_{-\frac{1}{2}[0,6,2;0,0]}$
$(000\underline{10}100)(0^8)'_{U_2}$	3	$(\mathbf{1},\mathbf{2})^L_{-\frac{1}{2}[0,12,0;0,0]}$
$(\frac{1}{2}\,\frac{1}{2}\,\frac{-1}{2}\,\underline{\frac{1}{2}\,\frac{-1}{2}}\,\frac{1}{2}\,\frac{1}{2}\,\frac{1}{2})(0^8)'_{U_3}$	2	$(\mathbf{1},\mathbf{2})^L_{-\frac{1}{2}[9,3,2;0,0]}$
$(0000\underline{0-1}0\underline{0-1})(0^8)'_{U_3}$	-4	$(\mathbf{1},\mathbf{2},\mathbf{2})^L_{\frac{1}{2}[0,-6,-2;0,0]}$
$(000\underline{\frac{2}{3}\,\frac{-1}{3}}\,\frac{2}{3}\,\frac{-1}{4}\,\frac{-1}{4})(\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{12}\,\frac{1}{12}\,\frac{1}{12})'_{T_{1_0}}$	0	$\star\ (\mathbf{1},\mathbf{2})^L_{\frac{1}{2}[0,6,-1;5,1]}$
$(000\underline{\frac{1}{3}\,\frac{-2}{3}}\,\frac{1}{3}\,\frac{1}{2}\,\frac{1}{2})(0^5\,\frac{-1}{3}\,\frac{-1}{3}\,\frac{-1}{3})'_{T_{2_0}}$	1	$(\mathbf{1},\mathbf{2})^L_{-\frac{1}{2}[0,0,2;0,-4]}$
$(0000\underline{0-1}0\,\frac{1}{4}\,\frac{1}{4})(\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{-1}{4}\,\frac{-1}{4}\,\frac{-1}{4})'_{T_3}$	-2	$(2_L+1_R)\cdot(\mathbf{1},\mathbf{2})^L_{\frac{1}{2}[0,-6,1;5,-3]}$
$(0000\underline{0-1}0\,\frac{-1}{4}\,\frac{-1}{4})(\frac{-1}{4}\,\frac{-1}{4}\,\frac{-1}{4}\,\frac{-1}{4}\,\frac{-1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4})'_{T_3}$	-4	$(2_L+1_R)\cdot(\mathbf{1},\mathbf{2})^L_{-\frac{1}{2}[0,-6,-1;-5,3]}$
$(000\underline{\frac{2}{3}\,\frac{-1}{3}}\,\frac{-1}{3}00)(0^5\,\underline{\frac{-2}{3}\,\frac{1}{3}\,\frac{1}{3}})'_{T_{4_0}}$	1	$6\cdot(\mathbf{1},\mathbf{2})^L_{\frac{1}{2}[0,0,0;0,0]}$
$(\frac{-1}{3}\,\frac{-1}{3}\,\frac{1}{3}\,\underline{\frac{2}{3}\,\frac{-1}{3}}\,\frac{2}{3}00)(0^8)'_{T_{4_-}}$	1	$\star\ 3\cdot(\mathbf{1},\mathbf{2})^L_{-\frac{1}{2}[-6,6,0;0,0]}$
$(\frac{1}{6}\,\frac{1}{6}\,\frac{-1}{6}\,\underline{\frac{1}{6}\,\frac{-5}{6}}\,\frac{1}{6}\,\frac{1}{2}\,\frac{1}{2})(0^8)'_{T_{4_-}}$	0	$2\cdot(\mathbf{1},\mathbf{2})^L_{-\frac{1}{2}[3,-3,2;0,0]}$
$(\frac{1}{6}\,\frac{1}{6}\,\frac{-1}{6}\,\underline{\frac{1}{6}\,\frac{-5}{6}}\,\frac{1}{6}\,\frac{-1}{2}\,\frac{-1}{2})(0^8)'_{T_{4_-}}$	-4	$2\cdot(\mathbf{1},\mathbf{2})^L_{\frac{1}{2}[3,-3,-2;0,0]}$
$(\frac{-1}{3}\,\frac{-1}{3}\,\frac{1}{3}\,\underline{\frac{1}{3}\,\frac{-2}{3}}\,\frac{1}{3}\,\frac{-1}{4}\,\frac{-1}{4})(\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{12}\,\frac{1}{12}\,\frac{1}{12})'_{T_{7_+}}$	-2	$(\mathbf{1},\mathbf{2})^L_{-\frac{1}{2}[-6,0,-1;5,1]}$
$(\frac{1}{6}\,\frac{1}{6}\,\frac{-1}{6}\,\underline{\frac{5}{6}\,\frac{-1}{6}}\,\frac{-1}{6}\,\frac{1}{4}\,\frac{1}{4})(\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{12}\,\frac{1}{12}\,\frac{1}{12})'_{T_{7_+}}$	3	$(\mathbf{1},\mathbf{2})^L_{-\frac{1}{2}[3,3,1;5,1]}$
$(000\underline{\frac{2}{3}\,\frac{-1}{3}}\,\frac{-1}{3}00)(0^5\,\underline{\frac{-2}{3}\,\frac{1}{3}\,\frac{1}{3}})'_{T_{4_0}}$	1	$3\cdot(\mathbf{1},\mathbf{2})^L_{\frac{-1}{2}[0,0,0;0,0]}$

Then, the $H_u - H_d$ type couplings arise from[3]

$$\epsilon_{\alpha\beta\gamma}\bar{\mathbf{3}}^{\alpha}_{W,I}\bar{\mathbf{3}}^{\beta}_{W,J}\bar{\mathbf{3}}^{\gamma}_{W,K}\epsilon^{IJK} \sim V\epsilon_{\alpha\beta}\bar{\mathbf{3}}^{\alpha}_{W,I}\bar{\mathbf{3}}^{\beta}_{W,J}\epsilon^{IJ}, \tag{23}$$

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, $\propto \bar{\mathbf{3}}_{W,\bar{1}}\bar{\mathbf{3}}_{W,\bar{2}} - \bar{\mathbf{3}}_{W,\bar{2}} \times \bar{\mathbf{3}}_{W,\bar{1}} = 0$. Because of ϵ^{IJ}, the same Higgs family does not have the coupling and the $H_u - H_d$ mass matrix of the 3×3 form 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 μ term [3,4].

3. Hidden sector $SU(5)'$, gauge mediation and messengers

As shown in Table 3, 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 $\mathbf{10}'$ and $\bar{\mathbf{5}}'$ is intact.

In Table 3, 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 $\mathbf{10}'$. Except $\mathbf{10}'$ of T_{2_-} and $\bar{\mathbf{5}}'$ of T_{7_0}, the rest 8 flavors form a vectorlike representation under $SU(5)'\times SU(2)_n \times U(1)_Y$. 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 $\mathbf{10}' + \bar{\mathbf{5}}'$ of a confining $SU(5)'$ breaks SUSY [11] and *we achieve the GMSB if the confining scale is below* 10^{12} GeV. Note that $\mathbf{10}'_0$ and $\bar{\mathbf{5}}'_0$ do not carry any $SU(3)_c \times SU(2)_W \times U(1)_Y$ charge (which is emphasized by the subscript 0) and DSB by $\mathbf{10}'_0$ and $\bar{\mathbf{5}}'_0$ does not break the SM gauge group.

Note that the singlet combination $\mathbf{10}'\mathbf{10}'\mathbf{10}'\bar{\mathbf{5}}'$ is not possible with one $\mathbf{10}'$. The $SU(5)'$ singlet combination in this uncalculable model can be parameterized by the gauge field strength field $\mathcal{W}'^{\alpha}\mathcal{W}'_{\alpha}$ as discussed in [13]. The interaction between the messenger f and the hidden sector gauge fields can appear from string compactification as

$$\mathcal{L} = \int d^2\theta \left[\xi(S_1, S_2, \ldots) f\bar{f}\mathcal{W}'^{\alpha}\mathcal{W}'_{\alpha} + \eta(S_1, S_2, \ldots) f\bar{f}\right] + \text{h.c.}, \tag{24}$$

where we have in general the holomorphic functions ξ and η of singlet chiral fields, $S_1, S_2, \ldots$. The quantum number of $\xi(S_1, S_2, \ldots) f\bar{f}$ 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 $\eta(S_1, S_2, \ldots) f\bar{f}$ should

[3] Note that $\mathbf{3}_W - \bar{\mathbf{3}}_W$ coupling is not generating $H_u - H_d$ terms since both H_u and H_d belong to $\bar{\mathbf{3}}_W$.

Table 3
Hidden sector $SU(5)'$ representations under $SU(2)_n \times SU(5)' \times SU(3)'$. After removing vectorlike representations by $\Gamma =$ even integer singlets, the starred representations remain

$P + n[V \pm a]$	Γ	No.$\times$(Repts.)$_{Y[Q_1,Q_2,Q_3,Q_4,Q_5]}$
$(\frac{1}{6}\,\frac{1}{6}\,\frac{-1}{6}\,\frac{1}{6}\,\frac{1}{6}\,\frac{1}{6}\,\frac{1}{4}\,\frac{1}{4})(\underline{\frac{-3}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}}\,\frac{-1}{4}\,\frac{-1}{4}\,\frac{-1}{4})'_{T1_-}$	2	$(\mathbf{1};\bar{\mathbf{5}}',\mathbf{1})^L_{0[3,3,1;1,-1]}$
$(\frac{1}{6}\,\frac{1}{6}\,\frac{-1}{6}\,\frac{-1}{6}\,\frac{-1}{6}\,\frac{-1}{6}\,00)(\underline{\frac{1}{2}\,\frac{1}{2}\,\frac{-1}{2}\,\frac{-1}{2}\,\frac{-1}{2}}\,\frac{-1}{6}\,\frac{-1}{6}\,\frac{-1}{6})'_{T2_+}$	-1	$\star(\mathbf{1};\mathbf{10}',\mathbf{1})^L_{0[3,-3,0;-2,-2]}$
$(0^6\,\underline{\frac{1}{4}\,\frac{-3}{4}})(\underline{\frac{3}{4}\,\frac{-1}{4}\,\frac{-1}{4}\,\frac{-1}{4}\,\frac{-1}{4}}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4})'_{T3}$	-1	$(\mathbf{2}_n;\mathbf{5}',\mathbf{1})^L_{0[0,0,-1;-1,3]}$
$(0^6\,\underline{\frac{3}{4}\,\frac{-1}{4}})(\underline{\frac{-3}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}}\,\frac{-1}{4}\,\frac{-1}{4}\,\frac{-1}{4})'_{T9}$	1	$(\mathbf{2}_n;\bar{\mathbf{5}}',\mathbf{1})^L_{0[0,0,1;1,-3]}$
$(0^3\,\frac{-1}{3}\,\frac{-1}{3}\,\frac{-1}{3}\,\frac{1}{4}\,\frac{1}{4})(\underline{\frac{-3}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}}\,\frac{1}{12}\,\frac{1}{12}\,\frac{1}{12})'_{T7_0}$	-1	$\star(\mathbf{1};\bar{\mathbf{5}}',\mathbf{1})^L_{0[0,-6,1;1,1]}$
$(\frac{1}{6}\,\frac{1}{6}\,\frac{-1}{6}\,\frac{1}{6}\,\frac{1}{6}\,\frac{1}{6}\,\frac{-1}{4}\,\frac{-1}{4})(\underline{\frac{3}{4}\,\frac{-1}{4}\,\frac{-1}{4}\,\frac{-1}{4}\,\frac{-1}{4}}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4})'_{T7_-}$	0	$(\mathbf{1};\mathbf{5}',\mathbf{1})^L_{0[3,3,-1;-1,3]}$
$(0^6\,\frac{-1}{2}\,\frac{-1}{2})(\underline{-10000}000)'_{T6}$	-2	$3\cdot(\mathbf{1};\bar{\mathbf{5}}',\mathbf{1})^L_{0[0,0,-2;-4,0]}$
$(0^6\,\frac{-1}{2}\,\frac{-1}{2})(\underline{10000}000)'_{T6}$	-2	$2\cdot(\mathbf{1};\mathbf{5}',\mathbf{1})^L_{1[0,0,-2;4,0]}$
$(0^6\,\frac{1}{2}\,\frac{1}{2})(\underline{-10000}000)'_{T6}$	2	$2\cdot(\mathbf{1};\bar{\mathbf{5}}',\mathbf{1})^L_{-1[0,0,2;-4,0]}$
$(0^6\,\frac{1}{2}\,\frac{1}{2})(\underline{10000}000)'_{T6}$	2	$3\cdot(\mathbf{1};\mathbf{5},\mathbf{1})^L_{0[0,0,2;4,0]}$

be $(-1,1,1)$. The H-momenta of the twisted sectors are given by [16,18,19]

$$U_1\colon\ (-1,0,0),\qquad U_2\colon\ (0,1,0),\qquad U_3\colon\ (0,0,1),$$
$$T_1\colon\ \left(\frac{-7}{12},\frac{4}{12},\frac{1}{12}\right),\qquad T_2\colon\ \left(\frac{-1}{6},\frac{4}{6},\frac{1}{6}\right),$$
$$T_3\colon\ \left(\frac{-3}{4},0,\frac{1}{4}\right),\qquad T_4\colon\ \left(\frac{-1}{3},\frac{1}{3},\frac{1}{3}\right),$$
$$T_5\colon\ \left(\frac{1}{12},\frac{-4}{12},\frac{-7}{12}\right),\qquad T_6\colon\ \left(\frac{-1}{2},0,\frac{1}{2}\right),$$
$$T_7\colon\ \left(\frac{-1}{12},\frac{4}{12},\frac{7}{12}\right),\qquad T_9\colon\ \left(\frac{-1}{4},0,\frac{3}{4}\right). \tag{25}$$

The Yukawa coupling $\eta(S_1,S_2,\ldots)f\bar{f}$ and the coefficient of $\mathcal{W}'^{\alpha}\mathcal{W}'_{\alpha}$ must satisfy the modular invariance rule for the twisted sector fields(z) multiplication,

$$\sum_z k(z)=0 \bmod 12,\qquad \sum_z [km_f](z)=0 \bmod 3. \tag{26}$$

Consider, for example, the vectorlike colored particles appearing only in $T4_-$ with $Q_{\rm em}=\mp\frac{1}{3}$: $f_3=D$, $\bar{f}_3=\bar{D}$, viz. Eq. (14). We can consider the following gauge singlet combination multiplied to $\mathcal{W}'^{\alpha}\mathcal{W}'_{\alpha}$, for example,

$$\begin{aligned} T4_-T4_-T1_0T7_+T4_-T4_- &\sim \bar{f}_3 f_3\langle T1_0T7_+T4_-T4_-\rangle \\ &\sim D_{-1/3}\bar{D}_{1/3}. \end{aligned} \tag{27}$$

Similarly, $SU(2)_W$ doublet coupling $\mathcal{W}'^{\alpha}\mathcal{W}''_{\alpha}$ can be considered. The product in (27), $T4_-T4_-T1_0T7_+T4_-T4_-$, 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, or f_1.

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

$$\mathcal{L}=\int d^2\theta\left[\frac{1}{M^2}f\bar{f}\mathcal{W}'^{\alpha}\mathcal{W}'_{\alpha}+M_f f\bar{f}\right]+\text{h.c.}, \tag{28}$$

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 ξ and η 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

$$M_{\rm mess}\approx M_f+\frac{\Lambda_h^3}{M^2},\qquad F_{\rm mess}\approx\frac{\Lambda_h^4}{M^2}. \tag{29}$$

With this GMSB scenario, firstly the observable sector gaugino obtains mass of order

$$\tilde{m}_{\rm SUSY}\sim\frac{g^2}{16\pi^2}\frac{\Lambda_h^4}{M^2M_{\rm mess}}, \tag{30}$$

while the gravitino mass is around $m_{3/2}\approx\Lambda_h^3/M_{Pl}^2$. 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 $(M^2M_{\rm mess})^{1/3}\approx 1.5\times 10^{14}$ GeV.

This leads us to consider the $\mathcal{W}'^{\alpha}\mathcal{W}'_{\alpha}$ couplings to H_uH_d and the observable sector Yukawa couplings $W_Y\sim H_uq_iu_j^c+H_dq_id_j^c$. Let us focus on the H_uH_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 $\int d^2\theta\, H_uH_d\mathcal{W}'^{\alpha}\mathcal{W}'_{\alpha}$ 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,

$$M_{\rm Higgsino}=\begin{pmatrix} 0, & A+a, & B+b \\ -A-a, & 0, & C+c \\ -B-b, & -C-c, & 0 \end{pmatrix}, \tag{31}$$

$$B_{\rm soft}=\mu\begin{pmatrix} 0, & a, & b \\ -a, & 0, & c \\ -b, & -c, & 0 \end{pmatrix}, \tag{32}$$

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 $\int d^2\theta\,(\frac{1}{m^3}H_u q u^c+\cdots)\mathcal{W}'^{\alpha}\mathcal{W}'_{\alpha}$ can be within a safe region of the gauge hierarchy solution. For example, the A-term estimated from this is $A\approx\frac{\Lambda_h^4}{m^3}$ 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 $\mathbf{10}'$ does not appear in the superpotential, it can appear in the Kähler potential. Possible terms of the form $\mathbf{10}'\mathbf{10}'^*ff^*/M_K^2$ might appear. The higher order Kähler terms was calculated for the compactification $T^6=(T_2)^3$ (with the volume moduli Ts and the complex structure moduli Cs) in Ref. [20] for two matter fields Q_α,

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

where n_α^i and $h_{2,1}=1$ (for our $\mathbf{Z}_{12-I}$) are the modular weight and a Hodge number, respectively. Also, l_α is an integer. The term $\mathbf{10}'\mathbf{10}'^*ff^*/M_K^2$ is not appearing in the above expression, and at present there does not exist a K^{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 $M^2/M_K^2<0.03$ [21].

4. Conclusion

We have shown that there exists a possibility of the hidden sector $SU(5)'$ with $\mathbf{10}'_0$ plus $\bar{\mathbf{5}}'_0$ matter below the GUT scale so that a GMSB at the stable vacuum is successful. Toward achieving the needed coupling constant α'_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 $U(1)_\Gamma$ gauge symmetry whose $\mathbf{Z}_2$ 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 $\frac{3}{8}$. 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.

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