SU(5)_{flip} \times SU(5)' from Z_{12-I}

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Abstract

Based on the Z_{12-I} orbifold compactification of the heterotic string theory, we construct a flipped-SU(5) model with three families of the standard model matter and ingredients for dynamical supersymmetry breaking. The doublet-triplet splittings in the Higgs representations 5_{-2} and 5_{2} are achieved by the couplings \[ 10_{1} \cdot 10_{1} \cdot 5_{-2} \text{ and } \bar{10}_{-1} \cdot \bar{10}_{-1} \cdot 5_{2}, \]
where 10_{1} and \bar{10}_{-1} develop GUT scale vacuum expectation values, breaking the flipped-SU(5) down to the standard model gauge group.
In this model, all the exotic states are decoupled from the low energy physics, and \( \sin^{2} \theta_{W}^{0} = \frac{3}{8} \).
Above the compactification scale, the flipped-SU(5) gauge symmetry is enhanced to the SO(10) gauge symmetry by including the Kaluza-Klein (KK) modes. The hidden sector gauge group is SU(5)'. The threshold correction by the KK modes allow a very wide range for the hidden sector confining scale (\( 10^{11} \text{ GeV} - 10^{16} \text{ GeV} \)). One family of hidden matter (\( \bar{10}' \) and \( 5' \)) gives rise to dynamical supersymmetry breaking.

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I. INTRODUCTION

The flipped-SU(5) \[\equiv SU(5)\times U(1)_X\] model, called SU(5)\textsubscript{flip}, was contrived for the alternative embedding of the standard model (SM) SU(2) singlets in the irreducible representations of the SU(5) grand unified theory (GUT) \[1, 2, 3, 25\] in contrast to the well-known Georgi-Glashow SU(5)\textsubscript{GG} \[4\]. As a result, a distinctive feature of the SU(5)\textsubscript{flip} is an interesting GUT breaking mechanism through the Higgs representation 10 of SU(5) rather than the adjoint 24, reducing the rank of the SU(5)\textsubscript{flip} by one unit. A great virtue of the SU(5)\textsubscript{flip} is the relative ease of the doublet/triplet splitting in the Higgs representations, 5 and 5, through a simple missing partner mechanism \[3\], which is also a result of such an embedding of the SM fermions. Another characteristic feature of the SU(5)\textsubscript{flip} is practically the absence of predicted fermion mass relations between quarks and leptons in contrast to the SU(5)\textsubscript{GG} GUT. As in the case of the SM, supersymmetric (SUSY) extension of the flipped-SU(5) \[2\] achieves also the gauge coupling unification with the LEP values of coupling constants \[5\], if the normalization of the hypercharge is assumed to be that of the SO(10) GUT, \(\sin^2\theta_W = \frac{3}{8}\).

With the advent of string constructions of the SUSY GUT models, and particularly, with the realization of the difficulty in obtaining adjoint Higgs for GUT breaking in string theory, the GUT breaking by the Higgs representations 10\(_1\) and 10\(_{-1}\) in the SU(5)\textsubscript{flip} became a great advantage. Earlier string construction obtaining 4-dimensional (4D) flipped-SU(5) GUTs was done in the fermionic construction \[6\]. Recently, a realistic model has been obtained in a \(\mathbb{Z}_{12-1}\) orbifold construction \[7\].

Since mid 1990s, dynamical SUSY breaking (DSB) toward phenomenological models has been advocated to resolve the SUSY flavor problem \[8\]. The well-known simple dynamical SUSY breaking (DSB) representations in the hidden sector are 10\(_{\prime}\) plus 5\(_{\prime}\) of SU(5)\textsubscript{\prime} \[9\], and 16\(_{\prime}\) of SO(10)\textsubscript{\prime} \[10\]. Other hidden sector gauge groups may be possible, but here we concentrate on a simple SU(5)\textsubscript{\prime} model with only one family, i.e. 10\(_{\prime}\) plus 5\(_{\prime}\), because of its relatively easy realization in heterotic string compactification. Recent DSB models at unstable vacua are known to be possible with vector-like representations in the hidden sector \[11\], which became popular because of our familiarity with SUSY QCD. For SU(N)\textsubscript{\prime}, the DSB requirement on the number of flavors in the SUSY QCD is \(N + 1 \leq N_f < \frac{3}{2}N\). For SU(5)\textsubscript{\prime}, \(N_f = 6\) and 7 satisfy this requirement. The DSB possibility from the heterotic string has been suggested by one of the authors at the unstable vacuum \[12\] and at the
stable vacuum \[13\]. In particular, the minimal supersymmetric standard model (MSSM) obtained in \[13\] with the SU(5)' gauge group and one family of \([10'] \oplus 5'\) in the hidden sector has many nice features such as the R-parity, one pair of Higgs doublets, and vector-like exotically charged states (exotics); but the bare value of the weak mixing angle is not \(\frac{3}{8}\).

The weak mixing angle would be, however, fitted to the observed one with the power-law type threshold effects contributed by the Kaluza-Klein towers \[14\], if relatively larger extra dimensions are assumed.

The so-called SUSY GUTs arise in two disguises: one is (usual) 4D SUSY GUTs such as the Dimopoulos-Georgi model \[15\] and the flipped-SU(5) \[2\], and the other GUTs in higher dimensions \((D > 4)\) as discussed in \[16\]. In a 4D SUSY GUT, the SM gauge group is obtained by spontaneous symmetry breaking of the GUT, whereas in a higher dimensional GUT it is achieved by the boundary conditions. String constructions of the MSSM \[14, 17, 19, 20, 21\] actually provided the idea of the higher dimensional SUSY GUT. In this paper, we will study a 4D SUSY GUT from a string compactification. In particular, based on the \(Z_{12-I}\) orbifold compactification of the heterotic string theory, we will construct a SUSY model SU(5)\(_{\text{flip}}\)×SU(5)', where the first (second) SU(5) indicates the gauge group of the visible (hidden) sector: The SU(5)\(_{\text{flip}}\) for the visible sector is broken to the SM gauge group by the vacuum expectation values (VEVs) of Higgs fields \(10'_H + \overline{10}'_H\), and the SU(5)' in the hidden sector becomes confined at lower energies, achieving DSB with one family of \([10' \oplus 5']\) \[9\].

This model yields MSSM fields plus one pair of \((10'_H + \overline{10}'_H)\) and one family of \([10' \oplus 5']\) in the hidden sector. All the other states in this construction are vector-like under the flipped-SU(5). A nice feature of the flipped-SU(5) model we present in this paper is that it gives a bare value \(\frac{3}{8}\) for \(\sin^2 \theta_W\). Above the compactification scale the visible sector flipped-SU(5) gauge symmetry is enhanced to SO(10) by including the KK modes.

This paper is organized as follows. In Secs. \(\text{II}\) and \(\text{III}\) we will construct a SUSY GUT model SU(5)\(_{\text{flip}}\)×SU(5)' and present the massless spectra from the untwisted and twist sectors. In Sec. \(\text{IV}\) we will discuss the Yukawa couplings needed for realization of the MSSM. Sec. \(\text{VI}\) will be devoted for the discussion of gauge coupling constants. Finally we will conclude in Sec. \(\text{VII}\).

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1 We use the one hidden sector matter notation as \([10' \oplus 5']\) of SU(5)' to distinguish it from the visible sector matter notation \([10' \oplus 5']\) of the flipped-SU(5).
II. $Z_{12-I}$ ORBITAL MODEL AND $U$ SECTOR FIELDS

We employ the $Z_{12-I}$ orbifold compactification scheme for the extra 6D space, which preserves $N = 1$ SUSY in the non-compact 4D spacetime [18, 19]. $Z_{12-I}$ orbifolds are known to give phenomenologically interesting MSSMs [7, 14, 19, 20].

The $Z_{12-I}$ orbifold is an SO(8) $\times$ SU(3) lattice, and the Wilson lines $W_3$ and $W_4 (= W_3)$ can be introduced in the 2D SU(3) lattice [18, 19]. We take the following shift vector $V$ and the Wilson line $W_3$,

$$V = \left( 0 0 0 0 0 ; -\frac{1}{6} -\frac{1}{6} -\frac{1}{6} \right) \left( 0 0 0 0 0 \frac{1}{4} \frac{1}{4} -\frac{1}{2} \right)'$$
$$W_3 = W_4 = W = \left( \frac{2}{3} \frac{2}{3} \frac{2}{3} \frac{2}{3} ; 0 -\frac{2}{3} -\frac{2}{3} \right) \left( \frac{2}{3} \frac{2}{3} \frac{2}{3} \frac{2}{3} 0 -\frac{2}{3} 0 0 \right)'$$

(1)

which are associated with the boundary conditions of the left moving bosonic string. For modular invariance in $Z_{12-I}$ orbifold compactification, $V$ and $W$ should be specially related to the twist vector $\phi = (\frac{5}{12} \frac{4}{12} \frac{1}{12})$, which is associated with the boundary conditions of the right moving superstrings, preserving only $N = 1$ SUSY in 4D. The twist vector $\phi = (\frac{5}{12} \frac{4}{12} \frac{1}{12})$ specifies the $Z_{12-I}$ orbifold. This model gives

$$V^2 - \phi^2 = \frac{1}{6}, \quad W^2 = \frac{16}{3}, \quad V \cdot W = -\frac{1}{6}. \quad (2)$$

Hence, the modular invariance conditions in $Z_{12-I}$ orbifold compactification are satisfied [19]:

- $12 \cdot (V^2 - \phi^2) =$ even integer,
- $12 \cdot W^2 =$ even integer,
- $12V \cdot W =$ integer.

The string excited states are irrelevant to low energy physics. The massless conditions for the left and right moving strings on the orbifold $Z_{12-I}$ are

left moving string:  \[
\frac{(P + kV_f)^2}{2} + \sum_i N_i^L \tilde{\phi}_i - \tilde{c}_k = 0,
\]

right moving string:  \[
\frac{(s + k\phi)^2}{2} + \sum_i N_i^R \tilde{\phi}_i - c_k = 0,
\]

where $k = 0, 1, 2, \cdots, 11$, $V_f = (V + m_f W)$ with $m_f = 0, +1, -1$, and $i$ runs over \{1, 2, 3, $\bar{1}, \bar{2}, \bar{3}$\}. Here $\tilde{\phi}_j \equiv k\phi_j$ mod $Z$ such that $0 < \tilde{\phi}_j \leq 1$, and $\tilde{\phi}_j \equiv -k\phi_j$ mod $Z$ such that $0 < \tilde{\phi}_j \leq 1$. $N_i^L$ and $N_i^R$ indicate oscillating numbers for the left and right movers. $P$ and $s \equiv (s_0, \bar{s})$ are the $E_8 \times E_8'$ and SO(8) weight vectors, respectively. The values of $\tilde{c}_k$, $c_k$ are found in Refs. [6, 18, 19].

The multiplicity for a given massless state is calculated with the GSO projector in the
The massless gauge sector corresponds to the states satisfying $P \cdot V = \text{integer}$, and $P \cdot W = \text{integer}$. They are

\begin{equation}
\text{SU}(5) : \quad (1 - 1 0 0 0 ; 0^3) (0^8)' \\
\text{SU}(5)' : \quad \begin{cases} (0^8) (1 - 1 0 0 0 0 0 0)' \\
\pm (0^8) (+- - - + + + +)' \\
\text{SU}(2)' : \quad \pm (0^8) (+ + + + + + + +)' ,
\end{cases}
\end{equation}

where the underline means all possible permutations. Thus, the gauge group is

\begin{equation}
[\text{SU}(5) \times \text{U}(1)_X] \times [\text{SU}(5) \times \text{SU}(2) \times \text{U}(1)^3]' ,
\end{equation}

where SU(5)×U(1)\text{X} is identified with the flipped SU(5). The U(1)_X charge operator of the flipped-SU(5) is

\begin{equation}
X = \frac{1}{\sqrt{40}} (-2 - 2 - 2 - 2 - 2 ; 0^3 ) (0^8)' .
\end{equation}

The normalization factor $\frac{1}{\sqrt{40}}$ is determined such that the norm of the $X$ (in general all U(1) charge operators in the level one heterotic string theory [19]) is $\frac{1}{\sqrt{2}}$. This value is exactly the one given as the normalization required for the SU(5)×U(1)_X embedded in SO(10). Since the standard model hypercharge is defined as

\begin{equation}
Y = \sqrt{\frac{3}{5}} (\frac{1}{3} \frac{1}{3} \frac{1}{3} \frac{1}{2} ; 0^3 ) (0^8)' ,
\end{equation}

\[ Z_{12-I} \text{ orbifold}, \]

\begin{equation*}
P_k(f) = \frac{1}{12 \cdot 3} \sum_{l=0}^{11} \tilde{\chi}(\theta^k, \theta^l) e^{2\pi i \Theta_k},
\end{equation*}

where $f (= \{f_0, f_+, f_-\})$ denotes twist sectors associated with $kV_f = kV, k(V+W), k(V-W)$. The phase $\Theta_k$ is given by

\begin{equation*}
\Theta_k = \sum_i (N_i^L - N_i^R) \hat{\phi}_i + (P + \frac{k}{2} V_f)V_f - (\tilde{s} + \frac{k}{2} \phi)\phi,
\end{equation*}

where $\hat{\phi}_j = \phi_j$ and $\hat{\phi}_j = -\phi_j$. Here, $\tilde{\chi}(\theta^k, \theta^l)$ is the degeneracy factor summarized in Ref. [6, 18, 19]. Note that $P_k(f_0) = P_k(f_+) = P_k(f_-)$ for $k = 0, 3, 6, 9$. In addition, the left moving states should satisfy

\[ P \cdot W = 0 \mod Z \] in the $U, T_3, T_6, T_9$ sectors.

The left moving states should satisfy

\[ P \cdot W = 0 \mod Z \] in the $U, T_3, T_6, T_9$ sectors.
Visible States | $P \cdot V$ | $\chi$ | SU(5)$_V$
---|---|---|---
(+ − − − ; + −)($0^8$)$_P$ | $\frac{1}{12}$ | L | 5$_3$
(+ + + − ; − +)($0^8$)$_P$ | $\frac{1}{12}$ | L | 10$_{-1}$
(+ + + + ; − + −)($0^8$)$_P$ | $\frac{1}{12}$ | L | 1$_{-5}$

TABLE I: The $U$ sector chiral states. There is no hidden sector chiral states and no flipped-SU(5) singlets.

the weak mixing angle at the string scale is $\sin^2 \theta_W = \frac{3}{8}$. From now on, we will drop the normalization factor “$\frac{1}{\sqrt{40}}$” and “$\sqrt{\frac{3}{5}}$” just for simplicity.

The massless chiral matter in the $U$ sector ($U$) are the states satisfying $P \cdot V = \{ \frac{-5}{12} \text{ or } \frac{4}{12} \text{ or } \frac{1}{12} \}$, and $P \cdot W = \text{integer}$. In Table I the chiral fields in the $U$ sector are tabulated. Note that there does not appear any flipped-SU(5) singlets in $U$. From the $U$ sector, we obtain one family of the MSSM matter

$$\overline{10}_{-1} + 5_3 + 1_{-5}, \ (\text{and their } CTP \text{ conjugates}),$$

where $\overline{10}_{-1}$, $5_3$, $1_{-5}$ contain $\{d_L^c, q_L, \nu_L^c\}$, $\{u_L^c, l_L\}$, and $e_L^c$, respectively. It is tempting to interpret this as the third (top quark) family, but the low dimensional Yukawa couplings prefer one in the twisted sector as the third family.

III. TWISTED SECTOR FIELDS

There are 11 twisted sectors, $T_k$ with $k = 1, 2, \cdots, 11$. The $CTP$ conjugates of the chiral states in $T_k$ is provided in $T_{12-k}$. Thus, it is sufficient to consider $k = 1, 2, \cdots, 6$. While the $U$ and $T_6$ sectors contain both chiral states and their $CTP$ conjugates, $T_1$, $T_2$, $T_3$, and $T_7$ ($T_{11}$, $T_{10}$, $T_8$, and $T_5$) sectors yield only the left-handed (right-handed) chiral states. The $T_3$ sector includes both left- and right-handed chiral states. So we will take $CTP$ conjugations for the right-handed states from the $T_3$ and $T_5$ sectors.
A. The flipped-$SU(5)$ spectrum

The visible sector chiral states of the twisted sectors are

$$T_4 : 2(\mathbf{10}_1 + \mathbf{5}_3 + \mathbf{1}_{-5}), \ 2(\mathbf{5}_{-2} + \mathbf{3}_2),$$  \hspace{1cm} (10)

$$T_3, T_9 : (\mathbf{10}_1 + \mathbf{10}_{-1}),$$  \hspace{1cm} (11)

$$T_7 : (\mathbf{5}_{-2} + \mathbf{3}_2),$$  \hspace{1cm} (12)

$$T_6 : 3(\mathbf{5}_{-2} + \mathbf{3}_2).$$  \hspace{1cm} (13)

To get the left-handed states from the $T_9$ and $T_7$ sectors, we acted the $\mathcal{CP}$ conjugations to the right-handed states of $T_3$ and $T_5$ sectors. From Table III (or Eq. (10)), we note that two families of the MSSM matter fields appear from $T_4$. Together with one family from the $U$ sector, thus, they form a three family model, including the three right-handed neutrinos.

| $P + 4V'$ | $x$ | $(N_L)_1$ | $P_4(f_0)$ | $SU(5)_X$ |
|------------|-----|-----------|-------------|------------|
| $\frac{1}{2} - \frac{1}{2} - \frac{1}{2} - \frac{1}{2}$ | L | 0 | 2 | $\mathbf{2} \cdot \mathbf{5}_3$ |
| $\frac{1}{2} + \frac{1}{2} + \frac{1}{2} - \frac{1}{2}$ | L | 0 | 2 | $\mathbf{2} \cdot \mathbf{1} - 3$ |
| $\frac{1}{2} + \frac{1}{2} + \frac{1}{2} + \frac{1}{2}$ | L | 0 | 2 | $\mathbf{2} \cdot \mathbf{5}_{-2}$ |
| $\frac{1}{2} + \frac{1}{2} + \frac{1}{2} + \frac{1}{2}$ | L | 0 | 2 | $\mathbf{2} \cdot \mathbf{3}_2$ |
| $0 0 0 0 0 ; \frac{1}{2} + \frac{1}{2} + \frac{1}{2} + \frac{1}{2}$ | L | 0 | 3 | $\mathbf{3} \cdot \mathbf{1}_0$ |
| $0 0 0 0 0 ; \frac{1}{2} + \frac{1}{2} + \frac{1}{2} + \frac{1}{2}$ | L | 1, 2, 13 | 2, 3, 2 | $(2 + 3 + 2) \cdot \mathbf{1}_0$ |
| $0 0 0 0 0 ; \frac{1}{2} + \frac{1}{2} + \frac{1}{2} + \frac{1}{2}$ | L | 1, 2, 13 | 2, 3, 2 | $(2 + 3 + 2) \cdot \mathbf{1}_0$ |
| $0 0 0 0 0 ; \frac{1}{2} + \frac{1}{2} + \frac{1}{2} + \frac{1}{2}$ | L | 1, 2, 13 | 2, 3, 2 | $(2 + 3 + 2) \cdot \mathbf{1}_0$ |

| $P + 4V''$ | $x$ | $(N_L)_1$ | $P_4(f_0)$ | $(SU(5)_X \cdot SU(3)_Y \cdot SU(2)_L)$ |
|------------|-----|-----------|-------------|--------------------------------|
| $\frac{3}{2} + \frac{1}{2} - \frac{1}{2} - \frac{1}{2}$ | L | 0 | 3 | $\mathbf{3} \cdot (1_{-5/3}; \mathbf{5}', \mathbf{1'})$ |
| $\frac{3}{2} + \frac{1}{2} + \frac{1}{2} - \frac{1}{2}$ | L | 0 | 2 | $\mathbf{2} \cdot (1_{-5/3}; \mathbf{1}', \mathbf{2'})$ |
| $\frac{3}{2} + \frac{1}{2} - \frac{1}{2} + \frac{1}{2}$ | L | 0 | 2 | $\mathbf{2} \cdot (1_{-5/3}; \mathbf{1}', \mathbf{2'})$ |
| $\frac{3}{2} + \frac{1}{2} + \frac{1}{2} + \frac{1}{2}$ | L | 0 | 2 | $\mathbf{2} \cdot (1_{-5/3}; \mathbf{1}', \mathbf{2'})$ |
| $\frac{3}{2} + \frac{1}{2} - \frac{1}{2} + \frac{1}{2}$ | L | 0 | 3 | $\mathbf{3} \cdot (1_{-5/3}; \mathbf{5}', \mathbf{1'})$ |
| $\frac{3}{2} + \frac{1}{2} + \frac{1}{2} - \frac{1}{2}$ | L | 0 | 2 | $\mathbf{2} \cdot (1_{-5/3}; \mathbf{1}', \mathbf{2'})$ |
| $\frac{3}{2} + \frac{1}{2} - \frac{1}{2} + \frac{1}{2}$ | L | 0 | 2 | $\mathbf{2} \cdot (1_{-5/3}; \mathbf{1}', \mathbf{2'})$ |

| $P + 4V_-$ | $x$ | $(N_L)_1$ | $P_4(f_0)$ | $(SU(5)_X \cdot SU(3)_Y \cdot SU(2)_L)$ |
|------------|-----|-----------|-------------|--------------------------------|
| $\frac{1}{2} - \frac{1}{2} - \frac{1}{2} - \frac{1}{2}$ | L | 0 | 3 | $\mathbf{3} \cdot (1_{-5/3}; \mathbf{5}', \mathbf{1'})$ |
| $\frac{1}{2} + \frac{1}{2} + \frac{1}{2} - \frac{1}{2}$ | L | 0 | 2 | $\mathbf{2} \cdot (1_{-5/3}; \mathbf{1}', \mathbf{2'})$ |
| $\frac{1}{2} - \frac{1}{2} - \frac{1}{2} + \frac{1}{2}$ | L | 0 | 2 | $\mathbf{2} \cdot (1_{-5/3}; \mathbf{1}', \mathbf{2'})$ |
| $\frac{1}{2} + \frac{1}{2} - \frac{1}{2} + \frac{1}{2}$ | L | 0 | 2 | $\mathbf{2} \cdot (1_{-5/3}; \mathbf{1}', \mathbf{2'})$ |

TABLE II: Chiral matter states in the $T_4^0$, $T_4^+$, and $T_4^-$ sectors. The multiplicities are shown as the coefficient in the last column.

In Table III some Higgs doublets are shown. Altogether, there appear six pairs of Higgs doublets from $T_4$, $T_7$ and $T_6$, among which therefore the candidates of the MSSM Higgs
doublets are chosen. We will explain in Sec. IV that except one pair of \( \{5_{-2}, \overline{5}_2\} \), the other pairs of five-plets with \( X = \pm 2 \) in the \( T_4 \), \( T_5 \), and \( T_6 \) sectors achieve superheavy masses, when some singlets under \([SU(5) \times U(1)_X] \times [SU(5) \times SU(2)]\) obtain VEVs of order the string scale. We regard the remaining one pair of \( \{5_{-2}, \overline{5}_2\} \) as the Higgs containing the MSSM Higgs. We will explain also how to decouple the triplets appearing in such five-plets in Sec. IV

| \( P \) | \( \chi \) | \( \Theta_{L,R} \) | \( T_a \) | \( SU(5)_X \) |
|---|---|---|---|---|
| \( (1, 0, 0, 0, 0) \) | \( (10, 0, 0, 0, 0) \) | \( \frac{1}{2} \) | \( 3 \) | \( 5_{-2} \) |
| \( (1, 0, 0, 0, 0) \) | \( (10, 0, 0, 0, 0) \) | \( \frac{1}{2} \) | \( 3 \) | \( \overline{5}_2 \) |
| \( (0, 0, 0, 0, 1) \) | \( (0, 0, 0, 0, 0) \) | \( \frac{1}{2} \) | \( 2 \) | \( 1_0 \) |
| \( (0, 0, 0, 0, 0) \) | \( (0, 0, 0, 0, 1) \) | \( \frac{1}{2} \) | \( 2 \) | \( 1_0 \) |
| \( (0, 0, 0, 0, 0) \) | \( (0, 0, 0, 0, 1) \) | \( \frac{1}{2} \) | \( 2 \) | \( 1_0 \) |

TABLE III: Massless states satisfying \( P \cdot W = 0 \) mod \( Z \) in \( T_6 \).

To break the flipped-SU(5) down to the SM, we need \( 10_1 \) (\( \equiv 10_H \)) and \( \overline{10}_{-1} \) (\( \equiv \overline{10}_H \)), which appear from \( T_3 \) and \( T_5 \) as shown in Table IV. As explained later, they couple to the \( \{5_{-2}, \overline{5}_2\} \) (\( \equiv \{5_h, \overline{5}_h\} \)) so that the pseudo-Goldstone mode \( \{D, D^c\} \) included in \( \{10_H, \overline{10}_H\} \) pair up with the triplets contained in \( \{5_{-2}, \overline{5}_2\} \) to be superheavy.

| \( P \) | \( \chi \) | \( \Theta_{L,R} \) | \( T_a \) | \( SU(5)_X \) |
|---|---|---|---|---|
| \( (1, 2, -1, 0, 0) \) | \( (10, 0, 0, 0, 0) \) | \( \frac{1}{2} \) | \( 1 \) | \( \overline{10}_1 \) |
| \( (1, 2, -1, 0, 0) \) | \( (10, 0, 0, 0, 0) \) | \( \frac{1}{2} \) | \( 1 \) | \( 10^* \), or \( (10_{-1} \text{ from } T_6) \) |
| \( (0, 0, 0, 0, 0) \) | \( (0, 0, 0, 0, 0) \) | \( \frac{1}{2} \) | \( 1 \) | \( 1 \) |
| \( (0, 0, 0, 0, 0) \) | \( (0, 0, 0, 0, 0) \) | \( \frac{1}{2} \) | \( 1 \) | \( 1_0 \) |
| \( (0, 0, 0, 0, 0) \) | \( (0, 0, 0, 0, 0) \) | \( \frac{1}{2} \) | \( 2 \) | \( 1_0^* \) |
| \( (0, 0, 0, 0, 0) \) | \( (0, 0, 0, 0, 0) \) | \( \frac{1}{2} \) | \( 2 \) | \( 1_0^* \) |
| \( (0, 0, 0, 0, 0) \) | \( (0, 0, 0, 0, 0) \) | \( \frac{1}{2} \) | \( 2 \) | \( 1_0^* \) |
| \( (0, 0, 0, 0, 0) \) | \( (0, 0, 0, 0, 0) \) | \( \frac{1}{2} \) | \( 2 \) | \( 1_0^* \) |

TABLE IV: Massless states from \( T_3 \). The starred chirality \( R \) states in \( T_3 \) can be represented also by un-starred chirality \( L \) states with the opposite quantum numbers in \( T_9 \).
The hidden-sector SU(5)′ spectrum

The hidden sector fields appear from twisted sectors. The chiral multiplets under SU(5)′×SU(2)′ are listed as follows.

\[ T_1 : 3(5', 1')_{-5/3}, \, 3(\overline{5}', 1')_{5/3}, \, 2(1', 2')_{-5/3}, \, 2(1', 2')_{5/3}, \]  
\[ T_2 : (1', 2')_{5/3}, \, (1', 2')_{-5/3}, \]  
\[ T_1 : (\overline{10}', 1')_0, \, (5', 2')_0, \, (\overline{5}', 1')_0, \, (1', 2')_0, \, (\overline{5}', 1')_{-5/3}, \, (1', 2')_{-5/3}, \, 2(1', 2')_{-5/3}, \]  
\[ T_7 : (5', 1')_{5/3}, \, 2(1', 2')_{-5/3}, \, (1', 2')_{5/3}. \]  

Here, we replaced again the right-handed states in the \( T_5 \) sector by the left-handed ones in \( T_7 \) by CT\( \mathcal{P} \) conjugations. We have not included non-abelian group singlets. The vectorlike representations in the above achieve superheavy masses when the neutral singlet under the flipped-SU(5) develop VEVs of order the string scale. We will discuss it in Sec. \[ \text{IV} \]

Removing vectorlike representations from Eqs. \[ (14) \] \[ (17) \], there remain

\[ (\overline{10}', 1')_0, \, (5', 2')_0, \, (\overline{5}', 1')_0, \, (1', 2')_0. \]  

The hidden sector SU(2)′ is broken by a GUT scale VEV of \( (1', 2')_0 \) of \[ (18) \]. Then, out of the representations of \[ (18) \], there remain one hidden sector family of SU(5)′

\[ \overline{10}'_0, \, 5'_0. \]  

which is the key toward the DSB with SU(5)′. \[ (19) \]

Representations in \[ (19) \] do not carry
any visible sector quantum numbers and the flipped-SU(5) is not broken by the DSB in the hidden sector. Our construction of one family SU(5)′ with $N_f = 0$ or 1 vector-like pair of 5′ and 5 does not change the fate of DSB due to the index theorem. But inclusion of supergravity effects gives a runaway solution at large values of the dilaton field $\Phi_{1/2}$. But the barrier separation between the SUSY breaking minimum and the runaway point must be very high. The barrier separation is controlled by the hidden sector scale.

Finally, in Table VI we list the so far neglected components of the vectorlike representations of the hidden sector fields carrying nonvanishing hypercharges.

TABLE VI: Chiral matter states satisfying $\Theta_{(0,\pm)} = 0$ in the $T_1^{(0,\pm)}$ sectors.
they could obtain superheavy masses, if the needed neutral singlets develop VEVs of order $\frac{1}{\Lambda}$. These are singlet exotics. Since they are also vector-like under the flipped SU(5), however, the $C\bar{T}\bar{P}$ conjugates with the left-handed chirality are provided by the states in the $T_7$ sector.

**C. The other vector-like exotic states**

The remaining charged states under the flipped-SU(5) are the singlets of SU(5)$\times$SU(5)$'$×SU(2)$'$. They are listed as follows.

\[
\begin{align*}
T_4 : & \quad 4 \cdot 1_{-5/3}, \ 4 \cdot 1_{5/3}, \\
T_2 : & \quad 1_{-10/3}, \ 2 \cdot 1_{5/3}, \ 1_{10/3}, \ 2 \cdot 1_{-5/3}, \\
T_1 : & \quad 1_{10/3}, \ 3 \cdot 1_{-5/3}, \ 1_{-10/3}, \ 2 \cdot 1_{5/3}, \\
T_7 : & \quad 1_{10/3}, \ 2 \cdot 1_{-5/3}, \ 1_{-10/3}, \ 3 \cdot 1_{5/3}.
\end{align*}
\]

These are singlet exotics. Since they are also vector-like under the flipped SU(5), however, they could obtain superheavy masses, if the needed neutral singlets develop VEVs of order the string scale. Hence, we can get the same low energy field spectrum as that of the MSSM. Such vector-like superheavy exotics could be utilized to explain the recently reported high energy cosmic positron excess.

| $P + 5V$ | $\chi$ (SU(5) X) | SU(5) X |
|----------|------------------|---------|
| $(0 \ 0 \ 0 \ 0 \ \frac{1}{2} \ \frac{1}{2} \ \frac{1}{2} \ \frac{1}{2}) \ (0 \ 0 \ 0 \ 0 \ \frac{1}{2} \ \frac{1}{2} \ \frac{1}{2} \ \frac{1}{2})$ | R 0 | 1 $\cdot$ 5' |
| $(1 \ 0 \ 0 \ 0 \ \frac{1}{2} \ \frac{1}{2} \ \frac{1}{2} \ \frac{1}{2}) \ (0 \ 0 \ 0 \ 0 \ \frac{1}{2} \ \frac{1}{2} \ \frac{1}{2} \ \frac{1}{2})$ | R 0 | 1 $\cdot$ 5' |
| $(0 \ 0 \ 0 \ 0 \ \frac{1}{2} \ \frac{1}{2} \ \frac{1}{2} \ \frac{1}{2}) \ (0 \ 0 \ 0 \ 0 \ \frac{1}{2} \ \frac{1}{2} \ \frac{1}{2} \ \frac{1}{2})$ | R 21 | 1 $\cdot$ 15 |
| $(0 \ 0 \ 0 \ 0 \ \frac{1}{2} \ \frac{1}{2} \ \frac{1}{2} \ \frac{1}{2}) \ (0 \ 0 \ 0 \ 0 \ \frac{1}{2} \ \frac{1}{2} \ \frac{1}{2} \ \frac{1}{2})$ | R 21 | 1 $\cdot$ 15 |

**TABLE VII:** Chiral matter states from the $T_5^{(0,\pm)}$ sectors. All of them are right-handed states.
IV. SINGLETS AND YUKAWA COUPLINGS

It is necessary to make exotics vectorlike and heavy. For this purpose, many singlets are required to develop large VEVs. In Table VIII we list singlet fields. At least, the following fields are given large VEVs at the string scale,

\[ S_2, S_3, S_4, S_5, S_7, S_{11}, S_{12}, S_{13}, S_{16}, S_{17}, S_{18}, S_{21}, S_{22}. \]  

(24)

These VEVs are possible through higher dimensional terms in the superpotential.

| indices | singlet states | \( x \) | \( P(f_0) \) | Label |
|---------|----------------|------|---------|-------|
| \( T_2 \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S1 |
| \( T_3 \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 11, 12, 13 | 2, 3, 2 | S2 |
| \( T_4 \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 11, 12, 13 | 2, 3, 2 | S3 |
| \( T_5 \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 11, 12, 13 | 2, 3, 2 | S4 |
| \( T_6 \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S5 |
| \( T_7 \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S6 |
| \( T_8 \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S7 |
| \( T_9 \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S8 |
| \( T_{10} \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S9 |
| \( T_{11} \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S10 |
| \( T_{12} \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S11 |
| \( T_{13} \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S12 |
| \( T_{14} \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S13 |
| \( T_{15} \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S14 |
| \( T_{16} \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S15 |
| \( T_{17} \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S16 |
| \( T_{18} \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S17 |
| \( T_{19} \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S18 |
| \( T_{20} \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S19 |
| \( T_{21} \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S20 |
| \( T_{22} \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S21 |
| \( T_{23} \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S22 |
| \( T_{24} \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S23 |
| \( T_{25} \) | \( \{0 0 0 0 0 ; \frac{1}{2} \frac{1}{2} \frac{1}{2} \frac{1}{2} \} \) | L | 0 | S24 |

TABLE VIII: Left-handed SU(5)×U(1)\(_X\)×SU(5)'×SU(2)' singlet states. The right-handed states in \( T_3 \) and \( T_5 \) are converted to the left-handed ones of \( T_9 \) and \( T_7 \), respectively.
A. Conditions

Neglecting the oscillator numbers, $H$-momenta of states in various sectors, $H_{\text{mom},0} \equiv (s + k\phi + r_-)$ are assigned as

\[
\begin{align*}
U_1 &: (-1,0,0), \quad U_2 : (0,1,0), \quad U_3 : (0,0,1), \\
T_1 &: (\frac{-7}{12},\frac{4}{12},\frac{1}{12}), \quad T_2 : (\frac{-1}{6},\frac{4}{6},\frac{1}{6}), \quad T_3 : (\frac{-3}{4},0,\frac{1}{4}), \\
T_4 &: (\frac{-1}{3},\frac{1}{3},\frac{1}{3}), \quad \{T_5 : (\frac{1}{12},\frac{-1}{12},\frac{-7}{12})\}, \quad T_6 : (\frac{1}{2},0,\frac{1}{2}), \\
T_7 &: (\frac{-1}{12},\frac{4}{12},\frac{7}{12}), \quad T_9 : (\frac{-1}{2},0,\frac{3}{2}), \\
& \quad (25)
\end{align*}
\]

from which $T_5$ will not be used since the chiral fields there are right-handed while the other fields are represented as left-handed. With oscillators, the $H$-momentum $\equiv (R_1,R_2,R_3)$ are

\[
(H_{\text{mom}})_j = (H_{\text{mom},0})_j - (N^L)_j + (N^L)_{\bar{j}}, \quad j = 1,2,3. \tag{26}
\]

The superpotential terms by vertex operators should respect the following selection rules [19]:

(a) Gauge invariance

(b) $H$-momentum conservation with $\phi = (\frac{5}{12},\frac{4}{12},\frac{1}{12})$,

\[
\sum_z R_1(z) = -1 \mod 12, \quad \sum_z R_2(z) = 1 \mod 3, \quad \sum_z R_3(z) = 1 \mod 12, \tag{27}
\]

where $z(\equiv A,B,C,\ldots)$ denotes the index of states participating in a vertex operator.

(c) Space group selection rules:

\[
\sum_z k(z) = 0 \mod 12, \quad \sum_z [km_f](z) = 0 \mod 3. \tag{28}
\]

If some singlets obtain string scale VEVs, however, the condition (b) can be merged into Eq. (28) in (c). Our strategy is to construct composite singlets (CSs) which have $H$-momenta, $(1\ 0\ 0), \ (-1\ 0\ 0), \ (0\ 1\ 0), \ (0\ -1\ 0), \ (0\ 0\ 1), \ (0\ 0\ -1)$, using only singlets developing VEVs of order at the string scale $M_{\text{string}}$. Then, with any integer set $(l\ m\ n)$, we can attach an appropriate number of CSs such that they make the total $H$-momentum $(-1\ 1\ 1)$. Since their VEVs are of order $M_{\text{string}}$, the Yukawa couplings multiplied by them are not suppressed.
B. Composite singlets

Specifically, let us consider a CS composed of $S_2$ with $(N^L)_j = 1_1$, $S_{21}$ with $(N^L)_j = \{2_3, 1_2\}$, and $S_{22}$ with $(N^L)_j = 2_1$ from $T^0_4$, $T^0_1$, and $T^0_7$, respectively. The CS, 

\[ S_2 S_{21} S_{22}, \]

fulfills the selection rules (a) and (c) and its H-momentum is calculated as

\[ \left[ \left( \frac{1}{3} \frac{1}{3} \frac{1}{3} \right) + (1 0 0) \right] + \left[ \left( \frac{7}{12} \frac{7}{12} \frac{7}{12} \right) + (0 -1 -2) \right] + \left[ \left( \frac{1}{12} \frac{1}{12} \frac{7}{12} \right) + (2 0 0) \right] = (2 0 -1). \]

The CS composed of $S_3$ with $(N^L)_j = 1_1$ or $1_2$ or $1_3$, $S_5 ((N^L)_j = 0)$, and $S_{17} ((N^L)_j = 2_3)$ from $T^0_4$, $T_6$, and $T^0_2$, respectively. 

\[ S_3 S_5 S_{17}, \]

fulfills also (a) and (c) and its H-momentum is given by $(0 1 -1), (-1 0 -1), \text{or} (-1 1 -2)$. Similarly, “$S_5 S_7$” satisfies “(a)” and “(c)” and gives the H-momentum of $(-1 0 1)$. By multiplying properly $S_2 S_{21} S_{22}, S_3 S_5 S_{17}, S_5 S_7$ (and their higher powers), thus, one can indeed construct CSs, whose H-momenta are $(1 0 0), (-1 0 0), (0 1 0), (0 -1 0), (0 0 1), (0 0 -1)$. For instance, $(1 0 0)$ can be obtained from $(2 0 -1) + (1 0 0)$, namely $(S_2 S_{21} S_{22})(S_5 S_7). (0 0 1)$ is achieved from $(S_2 S_{21} S_{22})(S_5 S_7)^2$.

Then we do not have to take care of the “H-momentum conservation” in the selection rule (b) for the superpotential. One can easily see that all the states in $T^+_4$ and $T^-_4$ achieve string scale masses by $\langle S_4 \rangle$. The states in \{T^+_2, T^-_2\}, \{T^+_1, T^-_1\} and \{T^+_1, T^-_7\} pair up to be superheavy by $\langle S_2 \rangle, \langle S_3 \rangle$, and $\langle S_4 \rangle$. Similarly, the singlet states in \{T^+_7, T^-_7\} and \{T^-_1, T^+_7\} pair up to be superheavy.

In order to break the flipped-SU(5) to the SM gauge group, we need GUT scale ($\approx$ string scale in our case) VEVs of $\overline{10}_H$ and $10_H$. We have them from $T_3$ and $T_9$, respectively. The term $10_H \overline{10}_H$ and terms with its higher powers are allowed. Thus, SUSY vacua where \[ \langle \overline{10}_H \rangle = \langle 10_H \rangle \approx M_{\text{string}} \approx M_{\text{GUT}} \] exist.

We regard a pair of $5_h$ and $\overline{5}_h$ in $T^+_4$ as the Higgs fields containing two Higgs doublets of the MSSM. For the missing partner mechanism, we need the couplings $\overline{10}_H \overline{10}_H 5_h$ and $10_H 10_H 5_h$. These couplings are allowed in the superpotential by multiplying CSs, $S_{18} S_{11} S_{16}$ and $S_{17} S_{12} S_{15}$, respectively.

The vector-like five-plets appearing in the $T^+_6$ sector obtain string scale masses. By $\langle S_21 \rangle$ one pair of five-plets in $T^0_7$ can pair up with one pair of five-plets in $T^0_4$ to be superheavy. The remaining one pair of the five-plets in $T^0_4$, i.e. \{5_h, $\overline{5}_h$\} can get a mass term (or $\mu$ term) by $S_{17} S_{18}$ and $S_1$. While a VEV $S_{17} S_{18}$ has been assumed, a VEV $S_1$ is not yet assumed. It can be determined by soft terms such that $\mu \equiv \langle S_{17} S_{18} + S_1 \rangle \approx m_{3/2}$ as in the next MSSM.

The MSSM matter states in the $T^0_4$ sector couple to the Higgs $5_h, \overline{5}_h$ in the same sec-
tor. Additionally $\langle S_2 S_3 S_4 \rangle$ can be multiplied to suppress the size of the Yukawa couplings. The matter states in the untwisted sector also can couple to them by $S_2$, $S_3$, and $S_4$: $\overline{10} - 1 \overline{10} - 1 5_h \times \langle S_2^2 \rangle$, $\overline{10} - 1 5_3 5_h \times \langle S_2 S_4 \rangle$, and $1 - 5_3 5_h \times \langle S_2 S_3 \rangle$. Since there are in total 21 $= (2 + 3 + 2) \times 3$ states in $S_2$, $S_3$, and $S_4$, they can be utilized to suppress the size of the Yukawa couplings.

C. White dwarf axions and one pair of Higgsino doublets

In this subsection, we comment how the needed horizontal symmetry can arise from our heterotic string compactification. But, we will not endeavor to discuss accidental global symmetries arising at some specific vacua [27, 28, 29]. In our previous paper [25], we introduced a variant very light axion to enhance the axion-electron coupling. This enhancement was motivated from the unexpected extra energy loss from the white dwarf evolution [30]. It is needed to distinguish families by the quantum numbers of an Abelian horizontal gauge symmetry $U(1)_H$ so that the mixing angles are of $O(10^{-1}) - O(10^{-3})$. The Peccei-Quinn symmetry broken at $\sim 10^{11}$ GeV cannot achieve this goal due to the small mixing $F_a/M_P \sim 10^{-7}$. Let us choose the $H$ direction as

$$H = \frac{1}{2} (1 1 1 1 1 3 - 1 1)(0 0 0 0 0 a b c)'$$

where

$$b = 2a - 20, \quad c = \frac{3}{2}a - 7.$$  \hspace{1cm} (30)

Then the $H$ quantum numbers of the visible sector quark and Higgs fields are shown below in the square brackets.

$$U : \overline{10} [0], \quad T_4 : 2 \overline{10} [0], \quad U : \overline{5} - 3 [0], \quad T_4 : 2 \overline{5} - 3 [-1],$$

$$T_4 : 2 \overline{5} - 2 [1], \quad T_4 : 2 5_2 [0], \quad T_7 : 5 - 2 [2], \quad 5_2 [1]$$

which has a $U(1)_H - SU(5)^2$ anomaly. But this anomaly is cancelled by the Green-Schwarz mechanism [31]. The $H$ quantum numbers of (32) are minus of those anticipated in Ref. [25], and hence can act as the needed horizontal gauge symmetry.

As seen in the previous subsection, one pair of quintet and anti-quintet in $T_7$ are coupled to one pair of quintet and anti-quintet in $T_4$ via $\langle S_{21} \rangle$, and the remaining the other pair in $T_4$ was assumed to contain the MSSM Higgs. In this subsection, we will assume that $\langle S_{21} \rangle$
and $\langle S_4 S_{16} \rangle$ is fine-tuned to be zero. It is possible because the quantum numbers of $S_{21}$ and $S_4 S_{16}$ are the same. Instead we need the following singlet VEVs to remove two pairs of Higgs quintet and anti-quintet,

$$T_4 : S_1 [-1] \text{ and/or } S_2 [-1],$$

$$T_1 : S_{19} [-2] \text{ and/or } S_{20} [-2].$$

The U(1)$_H$ invariant couplings of the form $T_1 T_4 T_4$ remove two pairs of Higgs quintet and anti-quintet of $T_4$. Note that in the previous subsection $\langle S_1 \rangle$ was adjusted to give a light mass mass term ("$\mu$ term") of one pair of the quintet and anti-quintet in $T_4$. The U(1)$_H$ invariant coupling of the form $T_1 T_4 T_7$ removes one pair of Higgs quintet and anti-quintet out of $T_4$ and $T_7$. Thus, the $3 \times 3$ Higgsino mass matrix takes the form,

$$
\begin{array}{ccc}
S_1[-1] & S_1[-1] & 0 \\
S_1[-1] & S_1[-1] & 0 \\
S_{19}[-2] & S_{19}[-2] & 0 \\
\end{array}
\begin{array}{c}
\bar{5}_2[0](T_4) \\
\bar{5}_2[0](T_4) \\
\bar{5}_2[1](T_7) \\
\end{array}
\begin{array}{c}
5^a_{-2}[1](T_4) \\
5^b_{-2}[1](T_4) \\
5^c_{-2}[2](T_7) \\
\end{array}
\text{(33)}
$$

It is obvious that $\bar{5}_2[1](T_7) \equiv \bar{5}^{EW}_{-2}$ is massless at this level. If $\langle S_1 \rangle = V_1$ and $\langle S_{19} \rangle = V_2$ and the Yukawa couplings are set to 1, the matching massless $5^{EW}_{-2}$ is a linear combination of fives from $T_4$ and $T_7$,

$$5^{EW}_{-2} = \frac{-V_2 (5^a + 5^b) + 2V_1 5^c}{\sqrt{4V_1^2 + 2V_2^2}}$$
\text{ (34)}

where the superscripts $a, b$ and $c$ denote their origins from $T_4$ and $T_7$ as indicated in Eq. \text{(33)}.

\section{V. KALUZA-KLEIN SPECTRUM}

The relatively light KK modes ($M_{KK} < 1/\sqrt{\alpha'}$) associated with the relatively large extra dimensions can arise only in the non-prime orbifolds such as the $Z_{12-I}$. It is because KK excitations are possible only under trivial (untwisted) boundary condition, which leads to $N = 2$ (or $N = 4$) SUSY spectra. In the $Z_{12-I}$ orbifold, for instance, the boundary conditions associated with the SU(3) sub-lattice of the 6D compact space in $U$, $T_3$, $T_6$, and $T_9$ sectors become trivial and allow $N = 2$ SUSY sectors.\text{[14].}
FIG. 1: The SU(3) lattice (a) and its dual lattice (b): (a) The torus is inside the yellow parallelogram and the fundamental region is the green parallelogram.

The KK modes associated with the relatively large extra dimensions \( R (\equiv R_3 = R_4) \) of the SU(3) sub-lattice, whose masses compose KK tower of \((\text{integer})/R\), should also satisfy the massless conditions \([14]\). Hence, the KK modes in the \( U \) sector still arise from the \( E_8 \times E_8' \) root vectors\(^2 \). But \( P \cdot W = \text{integer} \) is not necessary for the KK states in decompactification limit. In addition, the GSO projection condition in the \( U \) sector is relaxed from \( P \cdot V = \text{integer} \) to \( P \cdot 3V = \text{integer} \)[14]. The \( E_8 \times E_8' \) roots satisfying this are

\[
\begin{align*}
\text{SO}(10) & : \ (\pm 1 \pm 1 0 0 0 ; 0 0 0) (0^8)' \\
\text{SO}(6) & : \ (0 0 0 0 0 ; \pm 1 \pm 1 0) (0^8)' \\
E_6' & : \begin{cases} 
(0^8) (\pm 1 \pm 1 0 0 0 ; 0 0 0)' \\
(0^8) (+ - - - ; ++ +)' \\
(0^8) (+ + + + ; ++ +)' \\
(0^8) (+ + + + + ; ++ +)' 
\end{cases} \\
\text{SU}(2)_K' & : \ (0^8) (0 0 0 0 0 ; 1 - 1 0)' .
\end{align*}
\]

Thus, the gauge group is enhanced to

\[
[\text{SO}(10) \times \text{SO}(6)] \times [E_6 \times \text{SU}(2)_K \times \text{U}(1)]'.
\]

\(^2\) The states of \( E_8 \times E_8' \) weights, not satisfying \( P^2 = 2 \), are the string excited states with the masses of \((\text{integer})/\sqrt{\alpha'}\).
In the visible sector, the flipped-SU(5) in the massless case are embedded in a simple group SO(10). Therefore, between the GUT scale (≈ compactification scale) and string scale, the MSSM gauge couplings are unified in SO(10), including the U(1)_X coupling. SU(5)′ and SU(2)′ in the hidden sector are embedded in E′_6. Note that the SU(2)′_K emerging in 6D space is different from the SU(2)′ gauge symmetry observed from the massless spectrum. The SU(2)′ is embedded in the E′_6. The condition for KK matter states (N = 2 hypermultiplets) from the U sector is also relaxed from $P \cdot V = \{-\frac{5}{12}, \frac{4}{12}, \frac{1}{12}\} \pmod{Z}$ to $P \cdot 3V = \pm\frac{1}{4} \pmod{Z}$ [14]. The KK matter states from the U sector are shown in TABLE IX.

| Visible States | 4D χ | SO(10) × SO(6) |
|---------------|------|----------------|
| (\(\overline{16}\); + + −)(0^8)′ | L, R | (\(\overline{16}\), 4) |
| (\(\overline{16}\); + + +)(0^8)′ | L, R |

| Hidden States | 4D χ | E′_6 × SU(2)_K |
|---------------|------|----------------|
| (0^8)(\(\overline{16}\); + + −)′ | L, R | (27, 2)′ |
| (0^8)(±1 0 0 0 0 0 : 1 0 0)′ | L, R |
| (0^8)(0 0 0 0 0 ; −1 0 −1)′ | L, R |
| (0^8)(0 0 0 0 0 ; −1 0 1)′ | L, R | (1, 2)′ |

TABLE IX: The KK spectrum from the U sector. \(\overline{16}\) collectively denotes (+−−−−), (++++−−), and (+++−++), which are 5, \(\overline{10}\), and 1, respectively, in terms of SU(5). Here we drop the CTP conjugates.

Among the twisted sectors, only T_3, T_6 and T_9 can provide KK states in \(Z_{12−I}\). The KK states from T_9 are all the CTP conjugates of the KK states from T_3. As in the U sector, the KK modes from T_3, T_6, and T_9 should also satisfy the massless conditions. However, the required GSO projection is also relaxed. Following the guide of Ref. [14], one can derive the KK spectrum from the twisted sectors T_3 and T_6. The results are presented in TABLE X.

One can check that the KK spectra in TABLE IX and X cancel the 6D gauge anomalies. The beta function coefficients \(b^N_g\) of SO(10) and E′_6 by KK modes with \(N = 2\) SUSY are

\[
b^N_g = b^N_{\text{SO(10)}} = -2 \times 8 + 2 \times (2 \times 8 + 1 \times 10) = 36, \quad (40)
\]

\[
b^N_g = b^N_{\text{E}_6'} = -2 \times 12 + 2 \times 3 \times 2 = -12. \quad (41)
\]

The KK masses are nothing but the excited momenta (= \(\vec{m}_3, \vec{m}_4\)) in the SU(3) dual...
\[
\begin{array}{|c|c|c|c|}
\hline
P + 3V & T_k & (N^L)_j & 4D \chi \\
\hline
(0^5; +--)(0^5; \frac{3}{4} - \frac{1}{2})' & T_3 & 0 & L, R \\
(0^5; --)(0^5; \frac{3}{4} - \frac{1}{2})' & T_3 & 0 & L, R \\
(0^5; +--)(0^5; \frac{1}{4} + \frac{1}{2})' & T_3 & 11, 13 & L, R \\
(0^5; --)(0^5; \frac{1}{4} + \frac{1}{2})' & T_3 & 11, 13 & L, R \\
(16; 0 0 0)(0^5; --)(0^5; \frac{1}{4} + \frac{1}{2})' & T_3 & 0 & L, R \\
\hline
\end{array}
\]

\[4 \times (1, 4; 2')\]

\[8 \times (1, 4; 1')\]

\[4 \times (16, 1; 1')\]

\[3 \times (16, 1; 2')\]

\[5 \times (10, 1; 2')\]

\[
M_{KK}^2 = \frac{2g^{ab}}{3R^2} (m_a - P \cdot W) (m_b - P \cdot W)
\]

\[
M_{KK}^2 = \begin{cases} 
\frac{4}{3R^2} & \text{for } (m_3, m_4) = \pm(1, 0), \pm(0, 1), \pm(1, -1), \\
\frac{1}{R^2} & \text{for } (m_3, m_4) = \pm(1, 1), \pm(2, -1), \pm(1, -2).
\end{cases}
\]

\[
M_{KK}^2 = \begin{cases}
\frac{4}{9R^2} & \text{for } (m_3, m_4) = (0, 0), (1, 0), (0, 1), \\
\frac{16}{9R^2} & \text{for } (m_3, m_4) = (1, 1), (1, -1), (-1, 1).
\end{cases}
\]
FIG. 2: The KK modes with $P \cdot W = \text{integer}$. The length of the red arrows is $(4\alpha'/3L^2)^{1/2}$ and that of the blue arrows is $2(\alpha'/L^2)^{1/2}$.

FIG. 3: The KK modes with $P \cdot W = \frac{1}{3} \mod \text{integer}$. The lengths of the red, blue, and green arrows are $(\alpha'/L^2)^{1/2}$, $2(\alpha'/L^2)^{1/2}$, and $(7\alpha'/L^2)^{1/2}$, respectively.

In the next excited level, there are 6 KK states, whose mass-squareds are $\frac{28}{9R^2}$. The KK mass-squareds of the states with $P \cdot W = -\frac{1}{3} + \text{integer}$ and $(-m_3, -m_4)$ are the same as those of the states with $P \cdot W = \frac{1}{3} + \text{integer}$ and $(m_3, m_4)$. Non-vanishing vectors $(m_3, m_4)$ do not affect the GSO projection conditions. In TABLE XI, we display the KK states satisfying $P \cdot W = \text{integer}$. Except the states in TABLE XI thus, the other KK states in
The beta function coefficients $b_{SU(5)}^{N=2}$ by the “matter” states with $P \cdot W = \pm \frac{1}{3} + \text{integer}$ are

$$b_{SU(5)}^{N=2} = -2 \times 5 + 2 \times \left( \frac{1}{2} \times 12 + \frac{3}{2} \times 5 \right) = 17,$$

$$b_{U(1)_X}^{N=2} = \frac{1}{40} \times 2 \times (3^2 \times 10 + 1^2 \times 10 + 1^2 \times 40 + 2^2 \times 50) = 17,$$

$$b_{SU(5)'}^{N=2} = -2 \times 5 + 2 \times \frac{1}{2} \times 3 = -7.$$

The beta function coefficients $b_{b_{SU(5)}/4D}^{N=2}$ by the “CP conjugates” states with $P \cdot W = \pm \frac{1}{3} + \text{integer}$ are $b_{SU(5)}/4D^{N=2} - b_{SU(5)}^{N=2}$. Since $b_{SU(5)}^{N=2}$ is the same as $b_{U(1)_X}^{N=2}$, and both are included in $b_{SU(5)}^{N=2}$ in Eq. (10), the KK modes in this model do not affect the gauge coupling unification of SU(5) and U(1)$_X$. Accordingly, only the fields in $N = 1$ SUSY sector, which have no corresponding KK states, affect the unification.

TABLE XI: The KK spectrum satisfying $P \cdot W = \text{integer}$. Here we drop the CP conjugates.
From the beta function coefficients, we can expect that the MSSM gauge couplings rapidly increase in the ultraviolet region. On the other hand, the hidden sector gauge coupling is asymptotically free. Therefore, a large disparity in the visible and hidden sector couplings at the compactification scale can be unified to a single coupling at some scale above the compactification scale. It is interpreted as the string scale. In other words, starting with a unified coupling at string scale, the hidden sector SU(5)’ coupling can be of order one at a large scale.

When a gauge group $G$ is broken to a subgroup $H$ by Wilson line and further broken to $H_0$ by orbifolding ($H_0 = H$ in our model), the RG evolution of the gauge coupling of $H_0$, including the effects by KK modes, is described at low energies by

$$\frac{4\pi}{\alpha_{H_0}(\mu)} = \frac{4\pi}{\alpha_*} + b^{N=1}_{H_0} \log \frac{M_*^2}{\mu^2} + b^{N=2}_{H} \Delta^0 + b^{N=2}_{G/H} \Delta^\pm. \quad (49)$$

We assume that dilaton has been stabilized by a non-perturbative effect [32]. It can be discussed also in the context of SUSY breaking of Ref. [35]. In Eq. (49), $b^{N=2}_{H} \Delta^0$ ($b^{N=2}_{G/H} \Delta^\pm$) denotes the threshold correction by KK modes of $PW = 0 (\pm \frac{1}{3})$ mod integer, respecting $N = 2$ SUSY. $b^{N=2}_{H_0}$ in Eq. (49) is the beta function coefficient contributed by $N = 1$ SUSY sector states. As discussed above, the KK mass towers by the states with $P \cdot W = \frac{1}{3} +$ integer and with $P \cdot W = -\frac{1}{3} +$ integer are the same. $b^{N=2}_{G/H}$ is given by $b^{N=2}_{G} - b^{N=2}_{H}$.

As seen in Eqs. (40), (41), and (46), (47), (48), the beta function coefficients by KK modes are quite large. Accordingly, only the KK states residing in the lowest a few layers in the KK mass tower would be involved in the RG evolution of the visible SU(5) gauge coupling, before it reaches $O(1)$. So we will keep only such relatively light KK modes for RG analysis of the gauge couplings.

If $16/9R^2 < M_*^2 < 28/9R^2$, thus, $\Delta^0$ include the contributions by 6 KK modes with the mass-squared $4/3R^2$, while $\Delta^+ (and also \Delta^-)$ 3 KK modes of $4/9R^2$ and 3 of $16/9R^2$. Thus, the threshold corrections by such KK modes are given by

$$b^{N=2}_{H} \Delta^0 = 17 \cdot 6 \cdot \log \left( \frac{3R^2M_*^2}{4} \right), \quad (50)$$

$$b^{N=2}_{G/H} \Delta^\pm = 19 \cdot 3 \cdot 2 \left[ \log \left( \frac{9R^2M_*^2}{4} \right) + \log \left( \frac{9R^2M_*^2}{16} \right) \right], \quad (51)$$

where $H = SU(5)$ and $G = SO(10)$. We assume $1/R \approx M_{GUT}$ and $\alpha_* = 1$. With $\alpha_{SU(5)} = \frac{1}{25}$, we estimate $R^2M_*^2 \approx 2.5^3$, which is consistent with our assumption $16/9R^2 < M_*^2 < 28/9R^2$.

\footnote{Considering the first excited KK mass-squared is $4/9R^2$, one could define the effective compactification}
With $R^2 M_*^2 \approx 1.9$, and

$$b_{SU(5)'}^{N=1} = -3 \times 5 + \frac{1}{2} \times 3 + \frac{3}{2} = -12$$

(52)

by $(\mathbf{10}', \mathbf{1}')_0$, $(\mathbf{5}', \mathbf{2}')_0$, and $(\mathbf{5}', \mathbf{1}')_0$ in Eq. (19), one can estimate also the confining scale of the hidden SU(5)'. It is just below $\mu \approx 4/3R \approx 0.8 M_*$. Therefore, e.g. if $M_* = 2 \times 10^{16}$ GeV, the confining scale of the hidden sector is $1.6 \times 10^{16}$ GeV. Indeed, the string scale can be much lowered than $10^{18}$ GeV in the strongly coupled heterotic string theory (or the heterotic M theory), if the eleventh space dimension is sizable [33].

However, the hidden sector confining scale is very sensitive to $R^2 M_*^2$. If $M_*^2 \lesssim 4/9 R^2$, all the KK modes do not contribute to the RG evolution of the gauge couplings up to the string scale $M_*$, and so we should adopt only the usual 4D RG equation. If $M_* = 2 \times 10^{16}$ GeV and so $\alpha_{SU(5)'}^{-1} = 25$ at that scale, the confining scale can be much lower down to $10^{11}$ GeV. Here, we assumed SU(2)' is broken and only $\mathbf{10}'$ and $5'$ draw down the confining scale.

Below the confinement energy scale, the order parameters are composite fields rather than SU(5)’ gaugino and quarks. As noticed in Ref. [35], gaugino condensation scale or $N = 1$ SUSY breaking scale can be much lower than the confinement scale. Let us briefly discuss this issue in the following section.

VI. THE HIDDEN SECTOR SUPERSYMMETRY BREAKING

Now, let us proceed to consider the one family SU(5)’ model, with $\mathbf{10}'$ and $\mathbf{5}'$ plus $N_f$ copies of $5'$ and $\mathbf{5}'$. For $N_f = 0$ we can consider two composite chiral fields which are SU(5)’ singlets [13, 34],

$$W_\alpha^\alpha W_\beta^\beta,$$

$$\epsilon_{\alpha\gamma\eta\chi\xi} W_\alpha^\alpha W_\beta^\beta 10^{\alpha\beta} 5' 10^{\eta\delta} 10^{\chi\xi},$$

where $W_\alpha^\alpha$ is the hidden sector gluino superfield, satisfying $W_\alpha^\alpha = 0, (\alpha = 1, 2, \cdots , 5)$. There is no more SU(5)’-invariant independent chiral combination. For $N_f \neq 0$ also, due to the flavor symmetries of $5'$ and $\mathbf{5}'$, SU($N_f$) $\times$ SU($N_f + 1$), we consider only two composite SU(5)’ scale, $R_{\text{eff}} \equiv \frac{3}{2}R$. Then, $R_{\text{eff}}^2 M_*^2 = 5.6$. So at $\mu = M_*/\sqrt{5.6} = 0.4 \times M_*$, the first excited KK modes appear.
FIG. 4: A possible shape of the effective potential in terms of effective fields $Z$ and $Z'$. The lower curve is a schematic view including supergravity effects

singlet directions affected by instantons,

\[
Z \sim \mathcal{W}^{\rho}_{\beta} \mathcal{W}^{\beta}_{\alpha},
\]

\[
Z' \sim \epsilon^{\alpha\gamma\eta\xi} \mathcal{W}^{\rho}_{\beta} \mathcal{W}^{\gamma\delta}_{\delta} 10^{\nu\beta} 10^{\nu\delta} 10^{\nu\gamma}(5^{\nu\gamma\nu\mu}),
\]

where the lower indices $\nu$ and $\mu$ represent antisymmetric combinations. In terms of these composite chiral fields, it is known that the confining SUSY theory with one family is known to break SUSY dynamically. In this $F$-term breaking scenario, we can depict the SUSY breaking minimum as the local minimum in Fig. 4. In the lower curve, we show a schematic view including supergravity effects, which has a runaway piece at large values of $Z'$.

The dynamically generated effective superpotential, respecting these global symmetries plus the $5'$ flavor symmetry $SU(N_f)$ and the $\overline{5}'$ flavor symmetry $SU(N_f + 1)$, can be written as

\[
W_{SU(5)} = Z \left[ \log \left( \frac{Z^{2-N_f}Z'_{\Phi}}{\Lambda^{3N_c-2-N_f}} \right) - \alpha \right]
\]

where $\alpha$ is a coupling. It was shown that for $N_f = 3$, the SUSY conditions cannot be satisfied and SUSY is dynamically broken. Due to the index theorem, for any value of $N_f$, SUSY is dynamically broken, in particular in the $SU(5)'$ theory with one $10'$ and one $\overline{5}'$. The model with the fields of Secs. II and III has five flipped-$SU(5)$ families. But four of them carry the exotic $U(1)_X$ charges. So such four pairs should be assumed to be
superheavy to keep the gauge coupling unification. Nonetheless the model still contains the ingredients for the dynamical breaking of SUSY included. Inclusion of supergravity effects has been analyzed by one of us [35].

As discussed in Sec. V, the threshold correction by the KK modes allows a very wide range of the SU(5)' confinement scale, from $10^{11}$ GeV to $10^{16}$ GeV. Moreover, as noticed in Ref. [35], the gaugino condensation scale can be quite low compared to the confinement scale. Thus, even in the case where the confinement scale is above $10^{13}$ GeV, one can obtain $N = 1$ SUSY breaking effects in the visible sector of order $10^{2-3}$ GeV via the gravity mediation. If the condensation scale is below the $10^{13}$ GeV, SUSY breaking effects in the visible sector by the gauge mediation can dominate over those by the gravity mediation, and here one may resort to the gauge mediation scenario [13].

VII. CONCLUSION

We have constructed the flipped-SU(5)$\times$SU(5)' model with three families of the MSSM matter states, based on the $\mathbb{Z}_{12-I}$ orbifold compactification of the heterotic string theory. The flipped-SU(5) breaks down to the SM gauge group by non-zero VEVs of $10_H$ and $\overline{10}_H$. The doublet/triplet splitting problem is very easily resolved, because the missing partner mechanism simply works in flipped-SU(5). In this model, we could obtain $\sin^2\theta_W = \frac{3}{8}$ at the string (or GUT) scale as desired. We have shown that all the extra states beyond the MSSM field spectrum are vector-like under the flipped-SU(5) and obtain superheavy masses by VEVs of some neutral singlets.

In this model, the KK modes do not affect the gauge coupling unification in the visible sector, because the flipped-SU(5) gauge symmetry is enhanced to the SO(10) gauge symmetry above the compactification scale. On the other hand, they could cause a big difference between the visible and hidden gauge couplings at the compactification scale. Depending on the size of such disparity between the visible and hidden gauge couplings at the compactification scale, a wide range of the confining scale of SU(5)' is possible, $10^{11}$ GeV – $10^{16}$ GeV. With the hidden matter $\overline{10}'$ and $5'$, the gaugino condensation scale or the $N = 1$ SUSY breaking scale can be a few orders lower than the hidden sector SU(5)' confining scale.
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