Variations of the Hidden Sector in a Realistic Intersecting Brane Model

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Abstract

Recently, we discussed the first example of a phenomenologically realistic intersecting D6-brane model. In this model, the gauge symmetry in the hidden sector is \( USp(2)_1 \times USp(2)_2 \times USp(2)_3 \times USp(2)_4 \). However, we find that the \( USp(2)_1 \times USp(2)_2 \) gauge symmetry can be replaced by an \( U(2)_{12} \) gauge symmetry, and/or the \( USp(2)_3 \times USp(2)_4 \) gauge symmetry can be replaced by an \( U(2)_{34} \) gauge symmetry since the \( USp(2)^2 \) stacks of D6-branes contribute to the same Ramond-Ramond tadpoles as those of the \( U(2) \) stacks. Thus, there are three non-equivalent variations of the hidden sector, and the corresponding gauge symmetries are \( U(2)_{12} \times USp(2)_3 \times USp(2)_4 \), \( U(2)_{34} \times USp(2)_1 \times USp(2)_2 \), and \( U(2)_{12} \times U(2)_{34} \), respectively. Moreover, we study the hidden sector gauge symmetry breaking, discuss how to decouple the additional exotic particles, and briefly comment on the phenomenological consequences.

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I. INTRODUCTION

The goal of string phenomenology is to construct realistic standard-like string models with all moduli stabilized. In the early days, string model building was mainly concentrated on the weakly coupled heterotic string theory. After the second string revolution, consistent four-dimensional chiral models with non-Abelian gauge symmetry on Type II orientifolds were able to be constructed due to the advent of D-branes \[1\]. In particular, Type II orientifolds with intersecting D-branes, where the chiral fermions arise from the intersections of D-branes in the internal space \[2\] with T-dual description in terms of magnetized D-branes \[3\], have played an important role in string model building during the last few years.

On Type IIA orientifolds with intersecting D6-branes, many non-supersymmetric three-family standard-like models and Grand Unified Theories (GUTs) were constructed \[4, 5, 6\]. Although these models were globally consistent, there generically existed uncancelled Neveu-Schwarz-Neveu-Schwarz (NSNS) tadpoles as well as the gauge hierarchy problem. To solve these two problems, semi-realistic supersymmetric standard-like models, Pati-Salam models, SU(5) models as well as flipped SU(5) models have been constructed in Type IIA theory on \(T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)\) \[7, 8, 9, 10, 11, 12, 13, 14\] and \(T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2'\mathbb{Z}_2)\) \[15, 16\] orientifolds with intersecting D6-branes, and some of their phenomenological consequences have been studied \[17, 18\]. Moreover, the supersymmetric constructions in Type IIA theory on other orientifolds were also discussed \[19\]. There are two main constraints on supersymmetric D6-brane model building: RR tadpole cancellation conditions and four-dimensional \(N = 1\) supersymmetric D6-brane configurations. Also, K-theory conditions provide minor constraints. In addition, to stabilize the closed-string moduli via supergravity fluxes, the flux models on Type II orientifolds have also been constructed \[20, 21, 22, 23, 24, 25, 26, 27, 28, 29, 30\].

It is well known that there are two serious problems in almost all the supersymmetric D-brane models: no gauge coupling unification at the string scale, and the rank one problem in the Standard Model (SM) fermion Yukawa matrices. Although these problems can be solved in the flux models of Ref. \[29\] where the RR tadpole cancellation conditions are relaxed, these models are in the AdS vacua and the question of how to lift these AdS vacua to the Minkowski vacua or dS vacua correctly is still a big challenge. Recently, we found that there is one and only one intersecting D6-brane model on Type IIA \(T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)\) orientifold where the above problems can be solved \[11, 29\]. Moreover, this model may has a realistic low energy phenomenology \[31\]. Although its observable sector has unique phenomenological properties, it is possible to have different stacks of the D6-branes in the hidden sector.

In this paper, we discuss three non-equivalent variations of the hidden sector where the RR tadpoles are cancelled, the four-dimensional \(N = 1\) supersymmetry is perserved, and the K-theory conditions are satisfied. These three variations seem to be the only possibilities.
In the original model $[11, 29]$, the gauge symmetry in the hidden sector is $USp(2)_1 \times USp(2)_2 \times USp(2)_3 \times USp(2)_4$. Interestingly, we can replace the $USp(2)_1 \times USp(2)_2$ gauge symmetry by an $U(2)_{12}$ gauge symmetry, and/or the $USp(2)_3 \times USp(2)_4$ gauge symmetry by an $U(2)_{34}$ gauge symmetry since the contributions to the RR tadpoles from the $USp(2)^2$ stacks of D6-branes are the same as those of the $U(2)$ stacks. Thus, there are three non-equivalent variations, and the corresponding gauge symmetries in the hidden sector are $U(2)_{12} \times USp(2)_3 \times USp(2)_4$, $U(2)_{34} \times USp(2)_1 \times USp(2)_2$, and $U(2)_{12} \times U(2)_{34}$, respectively.

Moreover, we discuss the hidden sector gauge symmetry breaking, and consider how to decouple the additional exotic particles. Because the observable sector is the same, the discussions on phenomenological consequences, for example, the gauge coupling unification, supersymmetry breaking soft terms, low energy supersymmetric particle spectrum, dark matter density, and the SM fermion masses and mixings, are the same as those in Ref. $[31, 35]$.

This paper is organized as follows. We briefly review the intersecting D6-brane model building on Type IIA $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold in Section II and the realistic intersecting D6-brane model in Section III. We study the three variations of the hidden sector in Section IV. Discussion and conclusions are given in Section V.

II. INTERSECTING D6-BRANE MODEL BUILDING IN TYPE IIA THEORY ON $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ ORIENTIFOLD

We briefly review the intersecting D6-brane model building in Type IIA theory on $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold $[1, 3]$. We consider $T^6$ to be a six torus factorized as $T^6 = T^2 \times T^2 \times T^2$ whose complex coordinates are $z_i$, $i = 1, 2, 3$ for the $i$-th two torus, respectively. The $\theta$ and $\omega$ generators for the orbifold group $\mathbb{Z}_2 \times \mathbb{Z}_2$ act on the complex coordinates as following

$$\theta : (z_1, z_2, z_3) \rightarrow (-z_1, -z_2, z_3),$$

$$\omega : (z_1, z_2, z_3) \rightarrow (z_1, -z_2, -z_3).$$

We implement an orientifold projection $\Omega R$, where $\Omega$ is the world-sheet parity, and $R$ acts on the complex coordinates as

$$R : (z_1, z_2, z_3) \rightarrow (\bar{z}_1, \bar{z}_2, \bar{z}_3).$$

So, there are four kinds of orientifold 6-planes (O6-planes) for the actions of $\Omega R$, $\Omega R\theta$, $\Omega R\omega$, and $\Omega R\theta\omega$, respectively. Also, we have two kinds of complex structures consistent with orientifold projection for a two torus – rectangular and tilted $[32]$. If we denote the
TABLE I: General spectrum for intersecting D6-branes at generic angles, where $I_{aa'} = -2^{3-k} \prod_{i=1}^{3} (n^i_{a}l^i_{a})$, and $I_{aO6} = 2^{3-k} (-l^1_{a}l^2_{a}l^3_{a} + l^1_{a}n^2_{a}n^3_{a} + n^1_{a}l^2_{a}n^3_{a} + n^1_{a}n^2_{a}l^3_{a})$. Moreover, $\mathcal{M}$ is the multiplicity, and $a_S$ and $a_A$ denote the symmetric and anti-symmetric representations of $U(N_a/2)$, respectively.

| Sector     | Representation                                                                 |
|------------|-------------------------------------------------------------------------------|
| $aa$       | $U(N_a/2)$ vector multiplet and 3 adjoint chiral multiplets                    |
| $ab + ba$  | $\mathcal{M}(\frac{N_a}{2}, \frac{N_a}{2}) = I_{ab} = 2^{-k} \prod_{i=1}^{3} (n^i_{a}l^i_{b} - n^i_{b}l^i_{a})$ |
| $ab' + b'a$| $\mathcal{M}(\frac{N_a}{2}, \frac{N_a}{2}) = I_{ab'} = -2^{-k} \prod_{i=1}^{3} (n^i_{a}l^i_{b} + n^i_{b}l^i_{a})$ |
| $aa' + a'a$| $\mathcal{M}(a_S) = \frac{1}{2}(I_{aa'} - \frac{1}{2}I_{aO6})$ ; $\mathcal{M}(a_A) = \frac{1}{2}(I_{aa'} + \frac{1}{2}I_{aO6})$ |

For a stack of $N$ D6-branes that do not lie on the top of any O6-plane, we obtain the $U(N/2)$ gauge symmetry with three adjoint chiral superfields due to the orbifold projections. While for a stack of $N$ D6-branes on the top of an O6-plane, we obtain the $USp(N)$ gauge symmetry with three anti-symmetric chiral superfields. The bifundamental chiral superfields arise from the intersections of two different stacks of D6-branes or one stack of D6-branes and its $\Omega R$ image $\blacksquare \blacksquare$ In short, the general spectrum for intersecting D6-branes at generic angles, which is valid for both rectangular and tilted two tori, is given in Table I. Moreover, a model may contain additional non-chiral (vector-like) multiplet pairs from $ab + ba$, $ab' + b'a$, and $aa' + a'a$ sectors if two stacks of the corresponding D-branes are parallel and on the top of each other on one two torus. The multiplicity of the non-chiral multiplet pairs is given by the product of the intersection numbers on the other two two-tori.

Before further discussions, let us define the products of wrapping numbers

$$
A_a \equiv -n^1_{a}n^2_{a}n^3_{a}, \quad B_a \equiv n^1_{a}l^2_{a}l^3_{a}, \quad C_a \equiv l^1_{a}n^2_{a}l^3_{a}, \quad D_a \equiv l^1_{a}l^2_{a}n^3_{a},
$$

$$
\bar{A}_a \equiv -l^1_{a}l^2_{a}l^3_{a}, \quad \bar{B}_a \equiv l^1_{a}n^2_{a}l^3_{a}, \quad \bar{C}_a \equiv n^1_{a}l^2_{a}n^3_{a}, \quad \bar{D}_a \equiv n^1_{a}n^2_{a}l^3_{a}.
$$
TABLE II: Wrapping numbers of the four O6-planes.

| Orientifold Action | O6-Plane | (n₁, l₁) × (n₂, l₂) × (n₃, l₃) |
|--------------------|----------|-----------------------------|
| ΩR                | 1        | (2β₁, 0) × (2β₂, 0) × (2β₃, 0) |
| ΩRω               | 2        | (2β₁, 0) × (0, −2β₂) × (0, 2β₃) |
| ΩRθω              | 3        | (0, −2β₁) × (2β₂, 0) × (0, 2β₃) |
| ΩRθ               | 4        | (0, −2β₁) × (0, 2β₂) × (2β₃, 0) |

The four-dimensional \( N = 1 \) supersymmetric models from Type IIA orientifolds with intersecting D6-branes are mainly constrained by the RR tadpole cancellation conditions and the four-dimensional \( N = 1 \) supersymmetric D6-brane configurations, and also constrained by the K-theory conditions:

1. **RR Tadpole Cancellation Conditions**

The total RR charges of D6-branes and O6-planes must vanish since the RR field flux lines are conserved. And then we obtain the RR tadpole cancellation conditions as follows

\[
-2^k N^{(1)} + \sum_a N_a A_a = -2^k N^{(2)} + \sum_a N_a B_a = 0,
\]

\[
-2^k N^{(3)} + \sum_a N_a C_a = -2^k N^{(4)} + \sum_a N_a D_a = -16,
\]

where \( 2N^{(i)} \) are the number of D6-branes wrapping along the \( i \)-th O6-plane which is defined in Table II.

2. **Four-Dimensional \( N = 1 \) Supersymmetric D6-Brane Configurations**

The four-dimensional \( N = 1 \) supersymmetry can be preserved by the orientation projection if and only if the rotation angle of any D6-brane with respect to the O6-plane is an element of \( SU(3) \), or in other words, \( \theta_1 + \theta_2 + \theta_3 = 0 \mod 2\pi \), where \( \theta_i \) is the angle between the D6-brane and the O6-plane in the \( i \)-th two torus. This supersymmetry conditions can
be rewritten as

\[ x_A \tilde{A}_a + x_B \tilde{B}_a + x_C \tilde{C}_a + x_D \tilde{D}_a = 0, \]

\[ A_a/x_A + B_a/x_B + C_a/x_C + D_a/x_D < 0, \]  

(6)

where \( x_A = \lambda, \ x_B = \lambda 2^{\beta_2+\beta_3}/\chi_2 \chi_3, \ x_C = \lambda 2^{\beta_1+\beta_3}/\chi_1 \chi_3, \ x_D = \lambda 2^{\beta_1+\beta_2}/\chi_1 \chi_2, \) and \( \chi_i = R_i^2/R_1^1 \) are the complex structure parameters. The positive parameter \( \lambda \) has been introduced to put all the variables \( A, B, C, \) and \( D \) on an equal footing.

(3) K-theory Conditions

The discrete D-brane RR charges classified by the \( \mathbb{Z}_2 \) K-theory groups in the presence of orientifolds, which are subtle and invisible by the ordinary homology \([22, 33]\), should also be taken into account \([21]\). The K-theory conditions are

\[ \sum_a 2^{-k} \tilde{A}_a = \sum_a 2^{-\beta_1} N_a \tilde{B}_a = \sum_a 2^{-\beta_2} N_a \tilde{C}_a = \sum_a 2^{-\beta_3} N_a \tilde{D}_a = 0 \text{ mod } 4. \]  

(7)

III. THE REALISTIC INTERSECTING D6-BRANE MODEL

There may be one and only one intersecting D6-brane model in Type IIA theory on \( T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2) \) orientifold with a realistic phenomenology \([11, 29, 31]\). Let us briefly review it. We present the D6-brane configurations and intersection numbers in Table IIII and its spectrum in Table IVV. We put the \( a', b, \) and \( c \) stacks of D6-branes on the top of each other on the third two torus, and then we have the additional vector-like particles from \( N = 2 \) subsectors.

We have shown that the gauge symmetry in the observable sector can be broken down to the SM gauge symmetry via the Green-Schwarz mechanism, D6-brane splittings and supersymmetry preserving Higgs mechanism. The gauge couplings for \( SU(4)_C, SU(2)_L \) and \( SU(2)_R \) are unified at the string scale, and the additional exotic particles may be decoupled around the string scale. Also, we calculated the supersymmetry breaking soft terms, and the corresponding low energy supersymmetric particle spectrum that can be tested at the Large Hadron Collider (LHC). The observed dark matter density can also be generated. In addition, we can explain the SM quark masses and mixings, and the tau lepton mass. The neutrino masses and mixings may be generated via seesaw mechanism as well. Similar to the GUTs \([34]\), we have roughly the wrong fermion mass relation \( m_e/m_\mu \simeq m_d/m_s \), and the correct electron and muon masses can be generated via high-dimensional opera-
| stack | N   | (n_1,l_1)(n_2,l_2)(n_3,l_3) | A  | S   | b   | b'  | c   | c'  | O6^1 | O6^2 | O6^3 | O6^4 |
|-------|-----|-----------------------------|----|-----|-----|-----|-----|-----|------|------|------|------|
| a     | 8   | (0,-1)(1,1)(1,1)            | 0  | 0   | 30(3)| -3  | 0(3)| 1   | -1   | 0    | 0    |      |
| b     | 4   | (3,1)(1,0)(1,-1)            | -2 | 2   | -    | -   | 0(6)| 0   | 1    | 0    | -3   |      |
| c     | 4   | (3,-1)(0,1)(1,-1)           | 2  | -2  | -    | -   | -   | -1  | 0    | 3    | 0    |      |
| O6^1  | 2   | (1,0)(1,0)(2,0)             | -  | -   | -    | -   | -   | -   | -    | -    | -    | -    |
| O6^2  | 2   | (1,0)(0,-1)(0,2)            | -  | -   | -    | -   | -   | -   | -    | -    | -    | -    |
| O6^3  | 2   | (0,-1)(1,0)(0,2)            | -  | -   | -    | -   | -   | -   | -    | -    | -    | -    |
| O6^4  | 2   | (0,-1)(0,1)(2,0)            | -  | -   | -    | -   | -   | -   | -    | -    | -    | -    |

TABLE III: The D6-brane configurations and intersection numbers on Type IIA $\mathbb{T}^6/\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifold. The gauge symmetry is $[U(4)_C \times U(2)_L \times U(2)_R]_{\text{observable}} \times [U Sp(2)_1 \times U Sp(2)_2 \times U Sp(2)_3 \times U Sp(2)_4]_{\text{hidden}}$, the SM fermions and Higgs fields arise from the intersections on the first two torus, and the complex structure parameters are $2\chi_1 = 6\chi_2 = 3\chi_3 = 6$. Also, the beta functions for all $U Sp(2)_i$ gauge symmetries are $-3$.

Furthermore, all the $U Sp(2)_i$ gauge symmetries will become strong around the string scale [35].

IV. THREE VARIATIONS OF THE HIDDEN SECTOR

In the realistic intersecting D6-brane model [11,29], the observable sector is unique. Interestingly, we find three non-equivalent variations of the hidden sector where we can cancel the RR tadpoles, preserve the four-dimensional $N = 1$ supersymmetry, and satisfy the K-theory conditions. And it seems to us that there is no other variation. In the original model [11,29], the gauge symmetry in the hidden sector is $U Sp(2)_1 \times U Sp(2)_2 \times U Sp(2)_3 \times U Sp(2)_4$. We notice that the $U Sp(2)_1 \times U Sp(2)_2$ gauge symmetry can be replaced by an $U(2)_{12}$ gauge symmetry, and/or the $U Sp(2)_3 \times U Sp(2)_4$ gauge symmetry by an $U(2)_{34}$ gauge symmetry because the contributions to the RR tadpoles from the $U Sp(2)^2$ stacks of D6-branes are the same as those of the $U(2)$ stacks. Thus, there are three non-equivalent variations, and the corresponding gauge symmetries in the hidden sector are $U(2)_{12} \times U Sp(2)_3 \times U Sp(2)_4$, $U(2)_{34} \times U Sp(2)_1 \times U Sp(2)_2$, and $U(2)_{12} \times U(2)_{34}$, respectively. Let us present them one by one in the following subsections.
TABLE IV: The chiral and vector-like superfields, and their quantum numbers under the gauge symmetry $SU(4)_C \times SU(2)_L \times SU(2)_R \times USp(2)_1 \times USp(2)_2 \times USp(2)_3 \times USp(2)_4$.

| Field | Quantum Number | $Q_4$ | $Q_{2L}$ | $Q_{2R}$ |
|-------|----------------|-------|----------|----------|
| $ab$  | $3 \times (4, \overline{3}, 1, 1, 1, 1)$ | 1     | -1       | 0        |
| $ac$  | $3 \times (\overline{3}, 1, 2, 1, 1, 1)$ | -1    | 0        | 1        |
| $a1$  | $1 \times (4, 1, 1, 2, 1, 1, 1)$ | 1     | 0        | 0        |
| $a2$  | $1 \times (\overline{4}, 1, 1, 2, 1, 1)$ | -1    | 0        | 0        |
| $b2$  | $1 \times (1, 2, 1, 2, 1, 1)$ | 0     | 1        | 0        |
| $b4$  | $3 \times (\overline{3}, 1, 1, 1, 1, 2)$ | 0     | -1       | 0        |
| $c1$  | $1 \times (1, \overline{1}, 2, 1, 1, 1)$ | 0     | 0        | -1       |
| $c3$  | $3 \times (1, 1, 2, 1, 1, 2)$ | 0     | 0        | 1        |
| $b_S$ | $2 \times (1, 3, 1, 1, 1, 1)$ | 0     | 2        | 0        |
| $b_A$ | $2 \times (1, \overline{3}, 1, 1, 1, 1)$ | 0     | -2       | 0        |
| $c_S$ | $2 \times (1, \overline{1}, \overline{4}, 1, 1, 1)$ | 0     | 0        | -2       |
| $c_A$ | $2 \times (1, 1, 1, 1, 1, 1)$ | 0     | 0        | 2        |
| $a b'$ | $3 \times (4, 2, 1, 1, 1, 1)$ | 1     | 1        | 0        |
| $a c'$ | $3 \times (\overline{4}, 1, 2, 1, 1, 1)$ | -1    | -1       | 0        |
| $b c$ | $6 \times (1, 2, \overline{3}, 1, 1, 1)$ | 0     | 1        | -1       |
| $b c'$ | $6 \times (1, \overline{3}, 2, 1, 1, 1)$ | 0     | -1       | 1        |

A. $U(2)_{12} \times USp(2)_3 \times USp(2)_4$ Hidden Sector

In the first variation of the hidden sector, we replace the $USp(2)_1 \times USp(2)_2$ gauge symmetry by an $U(2)_{12}$ gauge symmetry. We present the D6-brane configurations and intersection numbers in Table [V]. Moreover, the particle spectrum has two parts: (1) the spectrum for old particles is given in Table [IV] by removing all the particles that are charged under $USp(2)_1 \times USp(2)_2$; (2) the spectrum for the new particles is given in Table [VI].

The anomalies from the global $U(1)$ of $U(2)_{12}$ are cancelled by the Green-Schwarz mechanism, and its gauge field obtains mass via the linear $B \wedge F$ couplings. Then, the effective gauge symmetry is $SU(2)_{12}$. The $SU(2)_{12}$ gauge symmetry can be broken down to $U(1)_{12}$ via D6-brane splitting. Interestingly, we do not have any additional chiral exotic particles that are charged under $SU(4)_C$. The simple way to give masses to the extra exotic particles would be by...
TABLE V: The D6-brane configurations and intersection numbers on Type IIA $T^6/Z_2 \times Z_2$ orientifold. The complete gauge symmetry is $[U(4)_C \times U(2)_L \times U(2)_R]_{\text{Observable}} \times [U(2)_{12} \times USp(2)_3 \times USp(2)_4]_{\text{Hidden}}$.

| Stack N | $(n_1,l_1)(n_2,l_2)(n_3,l_3)$ | A | S | $b$ | $b'$ | $c$ | $c'$ | d | $d'$ | $O6^3$ | $O6^4$ |
|---------|--------------------------------|---|---|-----|-----|-----|-----|---|-----|------|------|
| a       | 8 (0,-1) (1, 1) (1, 1)        | 0 | 0 | 3 0 | -3 0 | 0(2) | 0(1) | 0 | 0   |       |      |
| b       | 4 (3, 1) (1, 0) (1,-1)        | -2 | -2 | 0(6) | 0(1) | 1 0(1) | 0   | -3 |
| c       | 4 (3,-1) (0, 1) (1,-1)        | 2 2 | -2 2 | -1 0(1) | 3 0 |
| d       | 4 (1, 0) (1,-1) (1, 1)        | 0 0 | -2 -2 | - - | - - | -1 0(1) | 1 |
| $O6^3$  | 2 (0, -1) (1, 0) (0, 2)       | - - | - - | - - | - - | - - | - - |
| $O6^4$  | 2 (0, -1) (0, 1) (2, 0)       | - - | - - | - - | - - | - - | - - |

$X_{bd}$ and $X_{cd}$ is instanton effects [36, 37, 38, 39]. However, we do not have the suitable three-cycles wrapped by E2 instantons $^1$, and thus the instanton effects are not available. Similar results hold for the next two subsections. In addition, the $USp(2)_3$ and $USp(2)_4$ will become strong at about the string scale $^{[35]}$, and then we will have some composite particles in the $U(2)_{12}$ anti-symmetric and symmetric representations, $\overline{S}'_d$ and $T'_d$ from $X_{d3}$, and $S'_d$ and $T'_d$ from $X_{d4}$, respectively. So we can break the $U(1)_{12}$ by giving suitable string-scale vacuum expectation values (VEVs) to $\overline{T}'_d$ and $T'_d$, and we can give the string-scale VEVs to $\overline{S}'_d$ and $S'_d$. Note that we give the TeV-scale VEVs to $S'_L$ and the string-scale VEVs to $S'_R$ $^{[31]}$, we can give the GUT-scale masses to $X_{c3}^i$ and $X_{cd}$ and the TeV-scale masses to the $X_{bd}^i$ and $X_{cd}^i$.

$^1$ Note that the E2 branes must also wrap rigid cycles.
Furthermore, if we could give the string-scale masses to the three $U(2)_{12}$ adjoint chiral superfields and we do not break the $SU(2)_{12}$ via D6-brane splitting, the $SU(2)_{12}$ gauge symmetry will become strong around the string scale. Then we can have the singlet composite field $S_L'$ in the $U(2)_L$ anti-symmetric representation with charge +2 under $U(1)_L$ from $X_{bd}$. And we can give the string-scale VEVs to $S_i$ and $S_L'$ while keeping the D-flatness of $U(1)_L$. Therefore, we may also give the GUT-scale masses to the $X_{i4}$ and $X_{bd}$ via the high-dimensional operators [35].

B. $U(2)_{34} \times USp(2)_1 \times USp(2)_2$ Hidden Sector

| stack $N$ | $(n_1,l_1)(n_2,l_2)(n_3,l_3)$ | A | S | $b$ | $b'$ | c | $c'$ | e | $e'$ | $O6^1$ | $O6^2$ |
|---|---|---|---|---|---|---|---|---|---|---|---|
| a | 8 | (0,-1) | (1,1) | (1,1) | 0 | 0 | 3 | 0(3) | -3 | 0(3) | 0(2) | 0(0) | 1 | -1 |
| b | 4 | (3,1) | (1,0) | (1,-1) | -2 | -2 | - | 0(6) | 0(1) | 0(3) | -3 | 0 | 1 |
| c | 4 | (3,-1) | (0,1) | (1,-1) | 2 | -2 | - | - | - | - | 0(3) | 3 | -1 | 0 |
| e | 4 | (0,1) | (-1,1) | (-1,1) | 0 | 0 | - | - | - | - | - | - | - | 1 |
| $O6^1$ | 2 | (1,0) | (1,0) | (2,0) | - | - | - | - | - | - | - | - | - |
| $O6^2$ | 2 | (1,0) | (0,-1) | (0,2) | - | - | - | - | - | - | - | - | - |

TABLE VII: The D6-brane configurations and intersection numbers on Type IIA $T^6/Z_2 \times Z_2$ orientifold. The gauge symmetry is $[U(4)_C \times U(2)_L \times U(2)_R]_{\text{Observable}} \times [U(2)_{34} \times USp(2)_1 \times USp(2)_2]_{\text{Hidden}}$.

TABLE VIII: The new chiral superfields and their quantum numbers under the gauge symmetry $SU(4)_C \times SU(2)_L \times SU(2)_R \times U(2)_{34} \times USp(2)_1 \times USp(2)_2$. 

| Representation | Representation | $Q_4$ | $Q_{2L}$ | $Q_{2R}$ | $Q_{34}$ | Field |
|---|---|---|---|---|---|---|
| $bc'$ | $3 \times (1,\overline{1},\overline{1},1,1)$ | 0 | -1 | 0 | -1 | $X^i_{bc'}$ |
| $ce'$ | $3 \times (1,1,2,2,1,1)$ | 0 | 0 | 1 | 1 | $X^i_{ce'}$ |
| $e1$ | $1 \times (1,1,1,\overline{2},2,1)$ | 0 | 0 | 0 | -1 | $X_{e1}$ |
| $e2$ | $1 \times (1,1,1,2,1,2)$ | 0 | 0 | 0 | 1 | $X_{e2}$ |
In the second variation of the hidden sector, we replace the $USp(2)_3 \times USp(2)_4$ gauge symmetry by an $U(2)_{34}$ gauge symmetry. We present the D6-brane configurations and intersection numbers in Table VII. The particle spectrum also has two parts: (1) the spectrum for old particles is given in Table IV by removing all the particles that are charged under $USp(2)_3 \times USp(2)_4$; (2) the spectrum for the new particles is given in Table VIII.

Note that the wrapping numbers for the $d$ stack of D6-branes are equivalent to those of the $a$ stack by T duality and orientifold action, we can think that we have an $U(6)$ gauge symmetry in the beginning. Only the global $U(1)$ of $U(6)$ is anomalous $U(1)$ symmetry, and its gauge field obtains mass via the linear $B \wedge F$ couplings. After we put four D6-branes on the place with equivalent wrapping numbers (just like the D6-brane splittings), we break the $SU(6)$ down to the $SU(4)_C \times SU(2)_{34} \times U(1)'$ where the $U(1)'$ generator in $SU(6)$ is

$$T_{U(1)'} = \frac{1}{2\sqrt{6}} \text{diag}(1, 1, 1, 1, -2, -2).$$  

Thus, the left-handed and right-handed SM fermions have $U(1)'$ charges $+1/2\sqrt{6}$ and $-1/2\sqrt{6}$, respectively. In order to keep the gauge coupling unification, we have to break the $U(1)'$ so that it will not become part of the $U(1)_Y$. In short, we have to break $U(2)_{34}$ completely.

Because the $USp(2)_1$ and $USp(2)_2$ will become strong at about the string scale [35], we will have some composite particles in the $U(2)_{34}$ anti-symmetric and symmetric representations, $S'_e$ and $T'_e$ from $X_{e1}$, and $S'_e$ and $T'_e$ from $X_{e2}$, respectively. So we can break the $U(2)_{12}$ completely by giving suitable string-scale VEVs to $S'_e$, $T'_e$, $S'_e$, and $T'_e$. Moreover, we can have the singlet composite particle $S'_L$ in the $U(2)_L$ anti-symmetric representation with charge +2 under $U(1)_L$ from $X_{b2}$. And then we can give the string-scale VEVs to $S'_L$ and $S'_L$ while keeping the D-flatness of $U(1)_L$. Note that $S'_L$ also have string-scale VEVs, we may give the GUT-scale masses to $X_{b2}$, $X_{e1}$, $X_{be}$, and $X_{ce}$ via the high-dimensional operators [35]. Moreover, $X_{a1}$ and $X_{a2}$ may form the vector-like particles if we break the $USp(2)_1$ and $USp(2)_2$ down to the diagonal $USp(2)_{D12}$ [31].

C. $U(2)_{12} \times U(2)_{34}$ Hidden Sector

In the third variation of the hidden sector, we replace the $USp(2)_1 \times USp(2)_2$ gauge symmetry by $U(2)_{12}$, and replace the $USp(2)_3 \times USp(2)_4$ gauge symmetry by $U(2)_{34}$. We present the D6-brane configurations and intersection numbers in Table IX. The particle spectrum also has two parts: (1) the spectrum for old particles is given in Table IV by removing all the particles that are charged under $USp(2)_1 \times USp(2)_2 \times USp(2)_3 \times USp(2)_4$; (2) the spectrum for the new particles is given in Table X.
TABLE IX: The D6-brane configurations and intersection numbers on Type IIA $T^6/\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifold. The complete gauge symmetry is $[U(4)_C \times U(2)_L \times U(2)_R]_{\text{Observable}} \times [U(2)_{12} \times U(2)_{34}]_{\text{Hidden}}$.

TABLE X: The chiral and vector-like superfields, and their quantum numbers under the gauge symmetry $SU(4)_C \times SU(2)_L \times SU(2)_R \times U(2)_{12} \times U(2)_{34}$.

As discussed in above subsections, we can break the $U(2)_{12}$ down to the $U(1)_{12}$ gauge symmetry via Green-Schwarz mechanism and D6-brane splitting, and we have to break the $U(2)_{34}$ gauge symmetry completely. In order to break the $U(1)_{12}$ and $U(2)_{34}$ gauge symmetries, we put the $d$ and $e$ stacks of D6-branes on the top of each other on the second two torus, and put the $d$ and $e'$ stacks on the top of each other on the third two torus. Then, we have additional vector-like particles $X_{de}$, $X_{de'}$, $X_{de}$, and $X_{de'}$, as given in Table IX. And there exist the following Yukawa couplings

$$W \supset y^A_{ij} X_{bd} X_{be'}^i X_{de'}^j + y^B_{ij} X_{cd} X_{ce'}^i X_{de'}^j,$$

(9)

12
where $y^A_{ij}$ and $y^B_{ij}$ are Yukawa couplings. If we give the diagonal string-scale VEVs to $X_{de}^j$ and $X'_{de}$, we break the $U(2)_{12} \times U(2)_{34}$ down to the diagonal $U(2)_D$. Moreover, the $X_{bd}$ and one linear combination of $X'_{be}$, and the $X_{cd}$ and one linear combination of $X'_{ce}$, can have vector-like masses close to the string scale. Note that we can give the TeV-scale VEVs to $S_i^L$ and the string-scale VEVs to $S_i^R$ [31], we can give the GUT-scale masses to $X_{bd}$ and the other two linear combinations of $X'_{be}$, and the TeV-scale masses to $X_{cd}$ and the other two linear combinations of $X'_{ce}$, via the high-dimensional operators [35]. Similar to the discussions in the above subsection A, if we can give the string-scale masses to the three $U(2)_{12}$ adjoint chiral superfields and do not break the $SU(2)_{12}$ gauge symmetry via D6-brane splitting, the $SU(2)_{12}$ gauge symmetry will become strong around the string scale. Then we can have the singlet composite field $S'_L$ in the $U(2)_L$ anti-symmetric representation with charge $+2$ under $U(1)_L$ from $X_{bd}$, and we can give the string-scale VEVs to $S_i^L$ and $S'_L$ while keeping the D-flatness of $U(1)_L$. Therefore, we may also give the GUT-scale masses to $X_{bd}$ and the other two linear combinations of $X'_{be}$ via the high-dimensional operators [35].

V. DISCUSSION AND CONCLUSIONS

At present, there is only one known example of an intersecting D6-brane model with a realistic observable sector. Interestingly, there are three non-equivalent variations of the hidden sector in which the theoretical constraints on model building can be satisfied. There does not seem to be any other possible variation in the original model [11, 29], and the gauge symmetry in the hidden sector is $USp(2)_1 \times USp(2)_2 \times USp(2)_3 \times USp(2)_4$. We noticed that the $USp(2)_1 \times USp(2)_2$ gauge symmetry can be replaced by an $U(2)_{12}$ gauge symmetry, and/or the $USp(2)_3 \times USp(2)_4$ gauge symmetry can be replaced by an $U(2)_{34}$ gauge symmetry because the $USp(2)^2$ stacks of D6-branes contribute to the same RR tadpoles as those of the $U(2)$ stacks. Thus, we obtained three non-equivalent variations, and the corresponding gauge symmetries in the hidden sector are $U(2)_{12} \times USp(2)_3 \times USp(2)_4$, $U(2)_{34} \times USp(2)_1 \times USp(2)_2$, and $U(2)_{12} \times U(2)_{34}$, respectively. In addition, we studied the hidden sector gauge symmetry breaking, and discussed how to decouple the additional exotic particles. Because the observable sector is the same, the phenomenological discussions in the observable sector are the same as those in Ref. [31, 35].

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