Standard-Like Model Building on Type II Orientifolds

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(Dated: March 27, 2022)

Abstract

We construct new Standard-like models on Type II orientifolds. In Type IIA theory on $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold with intersecting D6-branes, we first construct a three-family trinification model where the $U(3)_C \times U(3)_L \times U(3)_R$ gauge symmetry can be broken down to the $SU(3)_C \times SU(2)_L \times U(1)_{Y_L} \times U(1)_{I_{LR}} \times U(1)_{I_{L}} \times U(1)_{Y_R}$ gauge symmetry by the Green-Schwarz mechanism and D6-brane splittings, and further down to the SM gauge symmetry at the TeV scale by Higgs mechanism. We also construct a Pati-Salam model where we may explain three-family SM fermion masses and mixings. Furthermore, we construct for the first time a Pati-Salam like model with $U(4)_C \times U(2)_L \times U(1)' \times U(1)''$ gauge symmetry where the $U(1)_{I_{LR}}$ comes from a linear combination of $U(1)$ gauge symmetries. In Type IIB theory on $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold with flux compactifications, we construct a new flux model with $U(4)_C \times U(2)_L \times U(2)_R$ gauge symmetry where the magnetized D9-branes with large negative D3-brane charges are introduced in the hidden sector. However, we can not construct the trinification model with supergravity fluxes because the three $SU(3)$ groups already contribute very large RR charges. The phenomenological consequences of these models are briefly discussed as well.

PACS numbers: 11.25.Mj, 11.25.Wx
I. INTRODUCTION

The major goal of string phenomenology is to construct four-dimensional models with the features of the Standard Model (SM), which connect the string theory to realistic particle physics. Due to the advent of D-branes, we can construct consistent four-dimensional chiral models with non-Abelian gauge symmetry on Type II orientifolds. Chiral matter can appear: (i) due to the D-branes located at orbifold singularities with chiral fermions appearing on the worldvolume of such D-branes; (ii) at the intersections of D-branes in the internal space with T-dual description in terms of magnetized D-branes.

On Type IIA orientifolds with intersecting D6-branes, a large number of non-supersymmetric three-family Standard-like models and grand unified models, which satisfy the Ramond-Ramond (RR) tadpole cancellation conditions, were constructed. However, there generically exist two problems: the uncancelled Neveu-Schwarz-Neveu-Schwarz (NSNS) tadpoles and the gauge hierarchy problem. On the other hand, the first supersymmetric models with quasi-realistic features of the supersymmetric Standard-like models have been constructed in Type IIA theory on $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold with intersecting D6-branes. Then, supersymmetric Standard-like models, Pati-Salam like models, $SU(5)$ models as well as flipped $SU(5)$ models have been constructed systematically, and their phenomenological consequences have been studied. In addition, supersymmetric constructions on other Type IIA orientifolds were also discussed.

In spite of these successes, the moduli stabilization in open string and closed string sectors is still an open problem even if some of the complex structure parameters (in the Type IIA picture) and dilaton fields may be stabilized due to the gaugino condensation in the hidden sector in some models (see, e.g., ). The supergravity fluxes give us another way to stabilize the compactification moduli fields by lifting continuous moduli space of the string vacua in an effective four-dimensional theory (see, e.g., ) because it introduces a supergravity potential. In this paper, for flux model building, we only consider the Type IIB orientifolds. The intersecting D6-brane constructions correspond to the models with magnetized D-branes where the role of the intersecting angles is played by the magnetic fluxes on the D-branes on the Type IIB orientifolds. In , the techniques for consistent chiral flux compactifications on Type IIB orientifolds were developed, and the dictionary for
the consistency and supersymmetry conditions between the two T-dual constructions was given. However, due to the Dirac quantization conditions, the supergravity fluxes impose strong constraints on consistent model building, since they contribute large positive D3-brane charges and then modify the global RR tadpole cancellation conditions significantly. Therefore, no explicit supersymmetric chiral Standard-like models were obtained in [25, 26]. In Type IIA orientifolds with flux compactifications, it has been recently shown that the RR, NSNS and metric fluxes could contribute negative D6-brane charges, and then relax the RR tadpole cancellation conditions [27, 28]. But we will not consider it in this paper.

By introducing the magnetized D9-branes with large negative D3-brane charges in the hidden sector, the first three-family and four-family Standard-like models with one unit of quantized flux were obtained [29, 30]. These constructions are T-dual to the supersymmetric models of intersecting D6-branes on $\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifold with the $Sp(2)_L \times Sp(2)_R$ or $Sp(2f)_L \times Sp(2f)_R$ gauge symmetry in the electroweak sector, respectively [18] (For the corresponding non-supersymmetric models, see [11]). Recently this kind of models has been studied systematically [31]. Moreover, considering the magnetized D9-branes with large negative D3-brane charges in the SM observable sector, a lot of new models have been constructed [32], for example, many new models with one unit of quantized flux, the first three- and four-family models with two units of quantized fluxes, and the first three- and four-family models with supersymmetric flux, i.e., three units of quantized fluxes.

However, in the previous model building with or without fluxes, one of the serious problems is how to generate suitable three-family SM fermion masses and mixings. In the $SU(5)$ models [16] and the flipped $SU(5)$ models [19], the up-type quark Yukawa couplings and the down-type quark Yukawa couplings are forbidden by anomalous $U(1)$ gauge symmetries. And in the Pati-Salam like models, although we can have the Yukawa couplings in principle, it is very difficult to construct three-family models which can give suitable SM fermion masses and mixings for three families because the left-handed fermions, the right-handed fermions and the Higgs fields in general arise from the intersections on different two-tori [33]. In addition, if supersymmetry is broken by supergravity fluxes, it seems that the masses for the rest massless SM fermions may not be generated from radiative corrections [33] because the supersymmetry breaking trilinear soft terms are universal and the supersymmetry breaking soft masses for the left/right-chiral squarks and sleptons are also universal [34, 35, 36] since the Kähler potential for the SM fermions depends only on the
intersection angles $37, 38, 39$. Thus, how to generate suitable three-family SM fermion masses and mixings is an interesting question. Another interesting question is whether we can construct new models with and without fluxes.

In this paper, in Type IIA theory on $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold with intersecting D6-branes, we first construct a three-family trinification model $[40, 41]$ with initial $U(3)_C \times U(3)_L \times U(3)_R$ gauge symmetry, which have not been studied previously (For the three-family trinification model building in the $Z_3$ orbifold compactification of weakly coupled heterotic string theory, see, e.g., $[42]$). Because of large RR charges from three $U(3)$ groups, it is very difficult to satisfy the RR tadpole cancellation conditions. In the trinification models, all the SM fermions and Higgs fields belong to the bi-fundamental representations, which can naturally arise from the intersecting D6-brane model building. Especially, the bi-fundamental representation with one fundamental and one anti-fundamental indices is different from the bi-fundamental representation with two fundamental (or anti-fundamental) indices, for example, $(3, \overline{3}, 1)$ and $(3, 3, 1)$ (or $(\overline{3}, \overline{3}, 1)$). So, the three-family trinification models can be constructed even if there is no tilted two-torus. Moreover, the $U(3)_C \times U(3)_L \times U(3)_R$ gauge symmetry is broken down to the $SU(3)_C \times SU(3)_L \times SU(3)_R$ gauge symmetry due to the Green-Schwarz (G-S) mechanism. Also, the $SU(3)_C \times SU(3)_L \times SU(3)_R$ gauge symmetry can be broken down to the $SU(3)_C \times SU(2)_L \times U(1)_{Y_L} \times U(1)_{Y_R}$ gauge symmetry by the splittings of the $U(3)_L$ and $U(3)_R$ stacks of the D6-branes $[18, 43]$, and the $U(1)_{Y_L} \times U(1)_{Y_R}$ gauge symmetry can be broken down to the $U(1)_{Y}$ gauge symmetry by the Higgs mechanism at the TeV scale which can not preserve the D-flatness and F-flatness and then breaks four-dimensional $N = 1$ supersymmetry because we assume the low scale supersymmetry in this paper. In addition, the quark Yukawa couplings are allowed while the lepton Yukawa couplings are forbidden by anomalous $U(1)$ gauge symmetries. In our model, we can only give masses to one family of the SM quarks.

For the Pati-Salam model, we consider the particles with quantum numbers $(4, \overline{2}, 1)$ and $(\overline{4}, 1, 2)$ under the $U(4)_C \times U(2)_L \times U(2)_R$ gauge symmetry as the SM fermions while we consider the particles with quantum numbers $(4, 2, 1)$ and $(\overline{4}, 1, \overline{2})$ as exotic particles because these particles are in fact distinguished by anomalous $U(1)$ gauge symmetries (In the trinification models, these kinds of particles are obviously different.). With this convention, we construct a three-family Pati-Salam model without tilted two-torus where the left-handed and right-handed SM fermions and the Higgs fields arise from the intersections on the same
two-torus, and then, we may explain three-family SM fermion masses and mixings. Also, the $U(4)_C \times U(2)_L \times U(2)_R$ gauge symmetry can be broken down to the $SU(3)_C \times SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$ gauge symmetry due to the G-S mechanism and D6-brane splittings. In our model, the $U(1)_{I_{3R}} \times U(1)_{B-L}$ gauge symmetry can only be broken down to the $U(1)_Y$ gauge symmetry at the TeV scale by Higgs mechanism.

In all the previous Pati-Salam like model building, the $U(1)_{I_{3R}}$ gauge symmetry arises from the non-Abelian gauge symmetry, for example, $U(2)_R$ or $USp(2f)_R$. We first construct a Pati-Salam like model with $U(4)_C \times U(2)_L \times U(1)' \times U(1)''$ gauge symmetry where the $U(1)_{I_{3R}}$ comes from a linear combination of $U(1)$ gauge symmetries. Also, the $U(1)_{I_{3R}} \times U(1)_{B-L}$ gauge symmetry can be broken down to the $U(1)_Y$ gauge symmetry at (or close to) the string scale by Higgs mechanism. However, only one family of the SM fermions can obtain masses.

In Type IIB theory on $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold with flux compactifications, we point out that one may not construct the trinification model with supergravity fluxes because the three $SU(3)$ groups already contribute very large RR charges. In addition, we construct a new flux model with $U(4)_C \times U(2)_L \times U(2)_R$ gauge symmetry where the magnetized D9-branes with large negative D3-brane charges are introduced in the hidden sector. This kind of models has not been studied previously because it is very difficult to have supersymmetric D-brane configurations with more than three stacks of $U(n)$ branes. However, in our model, the $U(1)_{I_{3R}} \times U(1)_{B-L}$ gauge symmetry can only be broken down to the $U(1)_Y$ gauge symmetry at the TeV scale by Higgs mechanism, and we can only give masses to one family of the SM fermions.

This paper is organized as follows. In Section II, we review the D-brane model building in Type II theories on the $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifolds. We consider the Type IIA non-flux model building and the Type IIB flux model building in Sections III and IV, respectively. Section V is our discussions and conclusions.

**II. INTERSECTING D-BRANE MODEL BUILDING**

In this Section, we briefly review the rules for the intersecting D-brane model building in Type II theories on $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifolds. It is well known that a supersymmetric Type IIA intersecting D6-brane construction is T-dual to a Type IIB construction.
with magnetized D3-, D5-, D7-, and D9-branes, thus, the rules for their model building are quite similar. For Type IIB constructions, we shall consider additional supergravity fluxes which give us stronger constraint on the RR tadpole cancellation conditions.

A. Type IIA Construction

We first briefly review the intersecting D6-brane model building in Type IIA theory on $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold [13, 14]. 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

\[
\theta : (z_1, z_2, z_3) \rightarrow (-z_1, -z_2, z_3),
\]
\[
\omega : (z_1, z_2, z_3) \rightarrow (z_1, -z_2, -z_3).
\] (1)

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

\[
R : (z_1, z_2, z_3) \rightarrow (\overline{z}_1, \overline{z}_2, \overline{z}_3).
\] (2)

With the wrapping numbers $(n^i, m^i)$ along the canonical bases of homology one-cycles $[a_i]$ and $[b_i]$, the general homology class for one cycle on the $i$-th two-torus $T^2_i$ is given by $n^i[a_i] + m^i[b_i]$. Therefore, the complete homology classes for the three cycles wrapped by a stack of $N_a$ D6-branes $[\Pi_a]$ and its orientifold image $[\Pi'_a]$ can be written as

\[
[\Pi_a] = \prod_{i=1}^3 (n^i_a[a_i] + m^i_a[b_i]), \quad [\Pi'_a] = \prod_{i=1}^3 (n^i_a[a_i] - m^i_a[b_i]).
\] (3)

The homology classes for four O6-planes associated with the orientifold projections $\Omega R$, $\Omega R\theta$, $\Omega R\omega$, and $\Omega R\theta\omega$ are

\[
\Omega R : [\Pi^{(1)}] = [a_1][a_2][a_3],
\]
\[
\Omega R\omega : [\Pi^{(2)}] = -[a_1][b_2][b_3],
\]
\[
\Omega R\theta : [\Pi^{(3)}] = -[b_1][a_2][b_3],
\]
\[
\Omega R\theta\omega : [\Pi^{(4)}] = -[b_1][b_2][a_3].
\] (4)
The four-dimensional $N = 1$ supersymmetric models from Type IIA orientifolds with intersecting D6-branes are mainly constrained in two aspects: RR tadpole cancellation conditions and four-dimensional $N = 1$ supersymmetry 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. With the filler branes on the top of the four O6-planes, we have the RR tadpole cancellation conditions:

\[ -N^{(1)} - \sum_a N_a n_a^1 n_a^2 n_a^3 = -N^{(2)} + \sum_a N_a n_a^1 m_a^2 m_a^3 = \]
\[ -N^{(3)} + \sum_a N_a m_a^1 n_a^2 m_a^3 = -N^{(4)} + \sum_a N_a m_a^1 m_a^2 n_a^3 = -16 , \]

where $N^{(i)}$ denotes the number of the filler branes on the top of the $i$-th O6-brane defined in Eq. (4).

2. **Four-Dimensional $N = 1$ Supersymmetry Conditions**

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 orientifold-plane is an element of $SU(3)$, i.e., $\theta_1 + \theta_2 + \theta_3 = 0 \mod 2\pi$, where $\theta_i$ is the angle between the D6-brane and the orientifold-plane on the $i$-th two-torus. This supersymmetry condition can be rewritten as

\[ -x_A m_a^1 m_a^2 m_a^3 + x_B m_a^1 n_a^2 n_a^3 + x_C n_a^1 m_a^2 n_a^3 + x_D n_a^1 n_a^2 m_a^3 = 0 , \]
\[ -n_a^1 n_a^2 n_a^3/x_A + n_a^1 m_a^2 m_a^3/x_B + m_a^1 n_a^2 m_a^3/x_C + m_a^1 m_a^2 n_a^3/x_D < 0 , \]

where $x_A = \lambda$, $x_B = \lambda^{2\beta_2+\beta_3}/\chi_2\chi_3$, $x_C = \lambda^{2\beta_1+\beta_3}/\chi_1\chi_3$, $x_D = \lambda^{2\beta_1+\beta_2}/\chi_1\chi_2$, $\chi_i = R_i^2/R_i^1$, are the complex structure parameters and $\lambda$ is a positive real number.

**B. Type IIB Construction**

We consider the Type IIB flux compactifications on the same orientifold $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ [25, 26]. The associated orientifold projection is still $\Omega R$, where the corresponding R acts on the complex coordinates as

\[ R : (z_1, z_2, z_3) \rightarrow (-z_1, -z_2, -z_3) . \]
We need $D(3+2n)$-branes to fill up the four-dimensional Minkowski space-time and wrap the $2n$-cycles on a compact manifold in Type IIB theory. And the introduction of magnetic fluxes provides the consistency to Type IIB theory. For one stack of $N_a$ D-branes wrapping $m^i_a$ times on the $i$-th two-torus $T^2_i$, we turn on $n^i_a$ units of magnetic fluxes $F_a$ for the center of mass $U(1)_a$ gauge factor on each $T^2_i$

$$m^i_a \frac{1}{2\pi} \int_{T^2_i} F^i_a = n^i_a .$$

(8)

Therefore, the D9-, D7-, D5- and D3-branes contain 0, 1, 2 and 3 vanishing $m^i_a$s, respectively.

The even homology classes for the point and the two-torus on the $i$-th two-torus $T^2_i$ are denoted as $[0_i]$ and $[T^2_i]$, respectively. Then, the vectors of RR charges of the $a$-th stack of D-branes and its $\Omega R$ image are

$$[\Pi_a] = \prod_{i=1}^{3} (n^i_a[0_i] + m^i_a[T^2_i]) , \quad [\Pi'_a] = \prod_{i=1}^{3} (n^i_a[0_i] - m^i_a[T^2_i]) ,$$

(9)

respectively. Similarly, the vectors of RR charges for the $O3$- and $O7_i$-planes, which respectively correspond to $\Omega R$, $\Omega R \omega$, $\Omega R \theta \omega$ and $\Omega R \theta$ O-planes, are

$$\Omega R : \quad [\Pi_{O3}] = [0_1] \times [0_2] \times [0_3] ;$$
$$\Omega R \omega : \quad [\Pi_{O7_1}] = -[0_1] \times [T^2_2] \times [T^2_3] ;$$
$$\Omega R \theta \omega : \quad [\Pi_{O7_2}] = -[T^2_1] \times [0_2] \times [T^2_3] ;$$
$$\Omega R \theta : \quad [\Pi_{O7_3}] = -[T^2_1] \times [T^2_2] \times [0_3] .$$

(10)

It is convenient to define the RR charges carried by the magnetized D-branes in Type IIB theory [32]. For one stack of $N_a$ D-brane with wrapping numbers $(n^i_a, m^i_a)$, the RR charges of D3-, D5-, D7-, and D9-branes are

$$Q_{3a} = N_a n^1_a n^2_a n^3_a , \quad Q_{5_{ia}} = N_a m^i_a m^j_a m^k_a ,$$
$$Q_{7_{ia}} = N_a n^i_a m^j_a m^k_a , \quad Q_{9_a} = N_a m^1_a m^2_a m^3_a ,$$

(11)

where $i \neq j \neq k$.

(1) RR Tadpole Concellation Conditions

The Type IIA RR tadpole cancellation conditions in Eq. [33] can also be applied exactly in Type IIB picture, except that there are flux contributions in Type IIB flux compactifications.
With the filler branes on the top of the $O3$- and $O7$-planes, we obtain the RR tadpole cancellation conditions

$$-N^{(O3)} - \sum_a Q3_a - \frac{1}{2} N_{flux} = -N^{(O7_1)} + \sum_a Q7_{1a} =$$

$$-N^{(O7_2)} + \sum_a Q7_{2a} = -N^{(O7_3)} + \sum_a Q7_{3a} = -16 , \tag{12}$$

where $N_{flux}$ is the amount of the fluxes turned on and quantized in units of elementary flux which is 64 as discussed later.

(2) Four-Dimensional $N = 1$ Supersymmetry Conditions

Four-Dimensional $N = 1$ supersymmetric vacua from flux compactifications require 1/4 supercharges of the ten-dimensional Type I theory be preserved in both open and closed string sectors. For the closed string sector, the specific Type IIB flux solution on orientifolds comprises of self-dual three-form field strength \[44, 45\]. The D3-brane RR charges contributed from the three-form flux $G_3 = F_3 - \tau H_3$ are given by

$$N_{flux} = \frac{1}{(4\pi^2 \alpha')^2} \int_{X_6} H_3 \wedge F_3 = \frac{i}{2 \Im(\tau)} \int_{X_6} G_3 \wedge G_3 , \tag{13}$$

where $F_3$ and $H_3$ are respectively the RR and NSNS three-form field strengths, and $\tau = \alpha + i/g_s$ is the Type IIB axion-dilaton coupling. Dirac quantization conditions for $F_3$ and $H_3$ on $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold require $N_{flux}$ to be a multiple of 64, and the BPS-like self-dual condition $\ast_6 G_3 = iG_3$ demands $N_{flux}$ to be positive. Supersymmetric configuration implies that $G_3$ background field should be a primitive self-dual (2,1) form, and a specific supersymmetric solution is \[45\]

$$G_3 = \frac{8}{\sqrt{3}} e^{-\pi i/6} (d\bar{z}_1 dz_2 dz_3 + dz_1 d\bar{z}_2 dz_3 + dz_1 dz_2 d\bar{z}_3) . \tag{14}$$

This fluxes stabilize the complex structure toroidal moduli at values

$$\tau_1 = \tau_2 = \tau_3 = \tau = e^{2\pi i/3} , \tag{15}$$

and its RR tadpole contribution $N_{flux}$ is 192.

In the open-string sector, the four-dimensional $N = 1$ supersymmetry for a D-brane configuration is conserved if and only if $\sum_i \theta_i = 0 \pmod{2\pi}$ is satisfied \[25\] where the “angle” $\theta_i$ is determined in terms of the worldvolume magnetic field as $\tan(\theta_i) \equiv (F_i)^{-1} = \frac{m_i^{1/2}}{n_i}$ and
\( \chi^i = R^i_1 R^i_2 \) is the area of the \( i \)-th two-torus \( T^2_1 \) in \( \alpha' \) units. This condition can be rewritten in terms of RR charges as 
\[ 
-x_A Q^9_a + x_B Q^5_{1a} + x_C Q^5_{2a} + x_D Q^5_{3a} = 0 ~, \\
-Q^3_a/x_A + Q^7_{1a}/x_B + Q^7_{2a}/x_C + Q^7_{3a}/x_D < 0 ~, 
\]  
(16)

where \( x_A = \lambda, x_B = \lambda/\chi'^2 \chi'^3, x_C = \lambda/\chi'^1 \chi'^3, x_D = \lambda/\chi'^1 \chi'^2 \).

It is not a surprise that the four-dimensional \( N = 1 \) supersymmetry conditions on Type IIB orientifold are similar to those on Type IIA orientifold in Eq. (6), except the different definitions of the complex structure parameters \( \chi^i \) and \( \chi'^i \). Therefore, we will take use of these two kinds of conditions on equal foot when we search the models.

C. Spectra

The spectra from the Type IIA and Type IIB orientifolds are the same. Massless chiral fields arise from the open strings with two ends attaching on the intersections of any two different D-brane stacks. The multiplicity (\( M \)) of the corresponding bi-fundamental representation is given by the intersection numbers between these two stacks of D-branes which is determined by the wedge product of their homology classes. The initial \( U(N_a) \) gauge group supported by a stack of \( N_a \) identical D-branes is broken down by the \( \mathbb{Z}_2 \times \mathbb{Z}_2 \) symmetry to a subgroup \( U(N_a/2) \) \[13, 14\]. The general spectra for the D-brane models on Type II orientifolds are given in Table I.

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 at least one two-torus. The multiplicity of the non-chiral multiplet pairs is given by the product of the rest intersection numbers, neglecting the null sector. For example, if only \( (n^i_{a1} l^i_{b1} - n^i_{b1} l^i_{a1}) = 0 \) in \( I_{ab} = [\Pi_a][\Pi_b] = \prod_{i=1}^3 (n^i_{a1} l^i_{b1} - n^i_{b1} l^i_{a1}) \), then we have
\[ 
\mathcal{M} \left( \left( \frac{N_a}{2}, \frac{N_b}{2} \right) + \left( \frac{N_a}{2}, \frac{N_b}{2} \right) \right) = \prod_{i=2}^3 (n^i_{a1} l^i_{b1} - n^i_{b1} l^i_{a1}) .
\]  
(17)

Moreover, the fermionic components of the non-chiral multiplets may acquire tree-level masses which dependent on the compactification radii and the brane wrapping numbers via Scherk-Schwarz mechanism \[46\].
| Sector | Representation |
|--------|----------------|
| $aa$   | $U(N_a/2)$ vector multiplet and 3 adjoint chiral multiplets |
| $ab + ba$ | $\mathcal{M}(N_a/2, N_b/2) = I_{ab} = \prod_{i=1}^{3}(n_i^a m_i^b - n_i^b m_i^a)$ |
| $ab' + b'a$ | $\mathcal{M}(N_a/2, N_b/2) = I_{ab'} = -\prod_{i=1}^{3}(n_i^a m_i^b + n_i^b m_i^a)$ |
| $aa' + a'a$ | $\mathcal{M}(\text{Anti}_{a}) = \frac{1}{2}(I_{aa'} + \frac{1}{2}I_{aO})$ |
|       | $\mathcal{M}(\text{Sym}_{a}) = \frac{1}{2}(I_{aa'} - \frac{1}{2}I_{aO})$ |

TABLE I: The general spectra for the D-brane models on Type II orientifolds, where $I_{aa'} = -8\prod_{i=1}^{3}n_i^a m_i^i$, and $I_{aO} = 8(-m_a^1 m_a^2 m_a^3 + m_a^2 n_a^2 n_a^3 + m_a^1 m_a^2 n_a^3 + n_a^1 n_a^2 m_a^3)$.

D. The K-theory Conditions

In addition to the RR tadpole cancellation conditions, the discrete D-brane RR charges classified by $\mathbb{Z}_2$ K-theory groups in the presence of orientifolds, which are subtle and invisible by the ordinary homology [29, 47], should also be taken into account [25].

In Type I superstring theory there exist non-BPS D-branes carrying non-trivial K-theory $\mathbb{Z}_2$ charges. To avoid this anomaly it is required that in compact spaces these non-BPS branes must exist in pairs [48]. Considering a Type I non-BPS D7-brane (\(\hat{D}7\)-brane), we know that it is regarded as a pair of D7-brane and its world-sheet parity image \(\overline{D}7\)-brane in Type IIB theory. Thus we require even numbers of these brane pairs in both Type IIA and IIB theories, since Type IIA theory is the T-dual of Type IIB theory. We only consider the effects from D3- and D7-branes since they do not contribute the standard RR charges which have been considered in the RR tadpole cancellation conditions. The K-theory conditions for a $\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifold, which are equivalent to the global cancellations of $\mathbb{Z}_2$ RR charges carried by the $\text{D5}_{a}^{i}$$\overline{\text{D5}}_{i}$ and $\text{D9}_{a}^{i}$$\overline{\text{D9}}_{i}$ brane pairs, were derived in [29] and are given by

$$\sum_{a} Q_{9a} = \sum_{a} Q_{51a} = \sum_{a} Q_{52a} = \sum_{a} Q_{53a} = 0 \text{ mod } 4 \ . \quad (18)$$
E. The Generalized Green-Schwarz Mechanism

Although the cubic non-Abelian anomalies in intersecting D-brane models are cancelled automatically when the RR tadpole cancellation conditions are satisfied, the additional mixed $U(1)$ anomalies may still be present. For instance, the mixed $U(1)$-gravitational anomalies which generate masses to the $U(1)$ gauge fields are not trivially zero [14, 20, 21]. These anomalies are cancelled by a generalized G-S mechanism which involves the untwisted RR forms. The couplings of the four untwisted RR forms $B_{i}^{j}$ to the $U(1)$ field strength $F_{a}$ of each stack $a$ are

$$
N_{a}m_{a}^{1}n_{a}^{2}n_{a}^{3}\int_{M_{4}}B_{2}^{1}\land trF_{a}^{*}, \quad N_{a}n_{a}^{1}m_{a}^{2}n_{a}^{3}\int_{M_{4}}B_{2}^{2}\land trF_{a}^{*},$

$$
N_{a}n_{a}^{1}n_{a}^{2}m_{a}^{3}\int_{M_{4}}B_{2}^{3}\land trF_{a}^{*}, \quad -N_{a}m_{a}^{1}m_{a}^{2}m_{a}^{3}\int_{M_{4}}B_{2}^{4}\land trF_{a}^{*} .
$$

These couplings determine the linear combinations of $U(1)$ gauge bosons that acquire string scale masses via the G-S mechanism. Sometimes, one combination of $U(1)$ gauge symmetries, for example $U(1)_{I_{3R}}$, must remain as an exact gauge symmetry because it is needed to generate the SM hypercharge interaction. Therefore, we must ensure that the $U(1)_{I_{3R}}$ gauge boson does not obtain such a mass. Suppose that the $U(1)_{I_{3R}}$ gauge symmetry is a linear combination of the $U(1)$s:

$$
U(1)_{I_{3R}} = \sum_{a}c_{a}U(1)_{a} .
$$

Because the corresponding field strength must be orthogonal to those that acquire masses via the G-S mechanism, we have

$$
\sum_{a}c_{a}Q5_{1a} = 0 , \quad \sum_{a}c_{a}Q5_{2a} = 0 ,
$$

$$
\sum_{a}c_{a}Q5_{3a} = 0 , \quad \sum_{a}c_{a}Q9_{a} = 0 .
$$

III. TYPE IIA MODEL BUILDING

In this Section, we shall construct the trinification model and the Pati-Salam like models on Type IIA orientifolds without flux.
A. Trinification Model

The $SU(3)_C \times SU(3)_L \times SU(3)_R$ trinification model, as a candidate for a grand unified theory, was proposed by de Rújula, Georgi, and Glashow \[40\] (see also \[41\]). Although no one has considered such models, the trinification model is quite interesting for the intersecting D-brane model building because all the left-handed quarks $Q^i_L$, the right-handed quarks $Q^i_R$, the leptons $L^i$, and the Higgs fields $H^k$, which are listed in Table II, belong to the bi-fundamental representations.

| Particles | Representation |
|-----------|----------------|
| $Q^i_L$   | $(3, \bar{3}, 1)$ |
| $Q^i_R$   | $(\bar{3}, 1, 3)$ |
| $L^i$ or $H^k$ | $(1, 3, \bar{3})$ |

TABLE II: The particle contents in the $SU(3)_C \times SU(3)_L \times SU(3)_R$ model.

Let us briefly review the trinification model. The electric charge generator $Q_{EM}$ is given by

$$Q_{EM} \equiv I_{3L} + \frac{Y}{2} = I_{3L} - \frac{Y_L}{2} + I_{3R} - \frac{Y_R}{2},$$

(22)

where the generators for $U(1)_{I_{3L}}$ and $U(1)_{I_{3R}}$, and $U(1)_{Y_L}$ and $U(1)_{Y_R}$ in $SU(3)_L$ and $SU(3)_R$ gauge symmetries are

$$T_{U(1)_{I_{3L}}R} = \begin{pmatrix} \frac{1}{2} & 0 & 0 \\ 0 & -\frac{1}{2} & 0 \\ 0 & 0 & 0 \end{pmatrix},$$

(23)

And the explicit particle components in the $(3, \bar{3}, 1)$, $(\bar{3}, 1, 3)$, and $(1, 3, \bar{3})$ representa-
The $SU(3)_C \times SU(3)_L \times SU(3)_R$ gauge symmetry can be broken down to the SM gauge symmetry by giving the vacuum expectation values (VEVs) to $\nu^c$ and $S$, i.e.,

$$\langle \nu^c \rangle \neq 0, \quad \langle S \rangle \neq 0.$$  

The electric charges for $h$ and $h^c$ are respectively $-\frac{1}{3}$ and $\frac{1}{3}$, for $E$ and $E^c$ are respectively $-1$ and $1$; and for $N$, $N^c$, and $S$ are zero.

With above background, we can construct an intersecting D6-brane trinification model. The bi-fundamental representation with one fundamental and one anti-fundamental indices is different from the bi-fundamental representation with two fundamental (or anti-fundamental) indices, for example, $(\bar{3}, 3, 1)$ and $(\bar{3}, 3, 1)$ (or $(3, \bar{3}, 1)$). So, we can construct the trinification models with three families of the SM fermions and without tilted two-torus.

There are three $SU(3)$ groups in the trinification model, so three stacks of six D6-branes are required. Additional stacks with $U(1)$ group and filler branes are also used to satisfy the RR tadpole cancellation conditions. In our model building, we require the intersection numbers to satisfy

$$I_{ab} = 3; \quad I_{ac} = -3; \quad I_{bc} \geq 4,$$  

where $I_{ab} = 3$ and $I_{ac} = -3$ give us three families of the left-handed quarks and three families of the right-handed quarks, respectively, and $I_{bc} \geq 4$ gives us three families of the leptons and $(I_{bc} - 3)$ Higgs field(s).

We have large RR charges from three $SU(3)$ groups, so it is not easy to construct a trinification model without RR tadpoles. After careful searches, we find a supersymmetric
intersecting D6-brane trinification model which satisfies the RR tadpole cancellation conditions and K-theory conditions. We present its complete wrapping numbers and intersection numbers in Table III and its spectrum in Table VII in Appendix A. In this model, we have three families of the SM fermions including the right-handed neutrinos, one pair of Higgs doublets, $H_u$ and $H_d$, one field $\nu^c$ and one field $S$.

### Table III: Wrapping numbers and intersection numbers in the $SU(3)_C \times SU(3)_L \times SU(3)_R$ model.

| stack | $N_a$ | $(n_1, l_1)(n_2, l_2)(n_3, l_3)$ | A  | b | b' | c  | c' | d  | d' | e  | e' | $f^{(1)}$ | $f^{(3)}$ |
|-------|------|---------------------------------|----|---|----|---|----|---|----|----|----|----------|----------|
| a     | 6    | (0, 1) (-1,-1) (2, 1)           | -2 | 2 | 3  | 1  | -3 | -3 | -3 | -3 | -3 | 2        | 0        |
| b     | 6    | (-1,-1) (-2, 1) (1, 0)          | -2 | 4 | 0(2)| 6  | 0(5)| 6  | 0(5)| 2  | -1       |          |
| c     | 6    | (-1, 1) (0, 1) (-1, 1)          | 0  | 0 | -  | -  | -  | 0(2)| 0(2)| 0(2)| 0  | 1        |          |
| d     | 2    | (-1, 1) (-1, 2) (1, 1)          | -16| 0 | -  | -  | -  | -  | 0(0)| -16| -1 | -2       |          |
| e     | 2    | (-1, 1) (-1, 2) (1, 1)          | -16| 0 | -  | -  | -  | -  | -  | -  | -1| -2       |          |
| fit(2)| 2    | (1, 0) (0, 1) (0,-1)            | -  | - | -  | -  | -  | -  | -  | -  | -  | -        | -        |
| fit(3)| 6    | (0, 1) (1, 0) (0,-1)            | -  | - | -  | -  | -  | -  | -  | -  | -  | -        | -        |

1) **Gauge Symmetry Breaking**

The $U(3)_C \times U(3)_L \times U(3)_R$ gauge symmetry is broken down to the $SU(3)_C \times SU(3)_L \times SU(3)_R$ gauge symmetry due to the G-S mechanism. And the $SU(3)_C \times SU(3)_L \times SU(3)_R$ gauge symmetry can be broken down to the $SU(3)_C \times SU(2)_L \times U(1)_{Y_L} \times U(1)_{I_{3R}} \times U(1)_{Y_R}$ gauge symmetry by the splittings of the $U(3)_L$ and $U(3)_R$ stacks of the D6-branes. Giving VEVs to the singlet Higgs fields $\nu^c$ and $S$, we can break the $U(1)_{Y_L} \times U(1)_{I_{3R}} \times U(1)_{Y_R}$ gauge symmetry down to the $U(1)_Y$ hypercharge interaction. The complete gauge symmetry breaking chains are

$$SU(3)_C \times SU(3)_L \times SU(3)_R$$

$$\xrightarrow{\text{Splitting}} SU(3)_C \times SU(2)_L \times U(1)_{Y_L} \times U(1)_{I_{3R}} \times U(1)_{Y_R}$$

$$\xrightarrow{\text{VEVs}} SU(3)_C \times SU(2)_L \times U(1)_Y .$$

We assume the low scale supersymmetry in this paper. Then the VEVs for $\nu^c$ and $S$ should be around the TeV scale because their Higgs mechanism can not preserve the D-flatness and F-flatness and then breaks four-dimensional $N = 1$ supersymmetry.

2) **Fermion Masses and Mixings**
The quark Yukawa couplings \( y_{ijk} Q^i_L Q^j_R H^k \) are allowed by the anomalous \( U(1) \) gauge symmetries in the intersecting D6-brane trinification model, while the lepton and neutrino Yukawa couplings \( y'_{ijk} L_i L_j H^k \) are forbidden by the anomalous \( U(1)_L \times U(1)_R \subset U(3)_L \times U(3)_R \) gauge symmetry.

In our model, only one family of the SM quarks can obtain masses because \( Q^i_L \) arise from the intersections on the second two-torus, while \( Q^i_R \) arise from the intersections on the third two-torus, or because we only have one pair of Higgs doublet fields.

### B. Pati-Salam Model

In the previous model building with or without fluxes, it is very difficult to generate suitable three-family SM fermion masses and mixings. In the \( SU(5) \) models and flipped \( SU(5) \) models, the up-type quark Yukawa couplings and the down-type quark Yukawa couplings are forbidden by anomalous \( U(1) \) gauge symmetries. And for the Pati-Salam like models, although all the Yukawa couplings could be allowed in principle, it is very difficult to construct three-family models which can give suitable masses and mixings to three families of the SM fermions because the left-handed fermions, the right-handed fermions and the Higgs fields in general arise from the intersections on different two-tori [33]. Moreover, if supersymmetry is broken by supergravity fluxes, it seems that the masses for the massless SM fermions may not be generated from radiative corrections [33] because the supersymmetry breaking trilinear soft terms are universal and the supersymmetry breaking soft masses for the left/right-chiral squarks and sleptons are also universal [34, 35, 36]. Thus, how to construct the Standard-like models, which can give suitable fermion masses and mixings for three families, is an interesting problem.

To solve this problem, we construct another class of supersymmetric Pati-Salam models without RR tadpoles and K-theory anomaly. In particular, under the \( U(4)_C \times U(2)_L \times U(2)_R \) gauge symmetry, we consider the particles with quantum numbers \((4, \overline{2}, 1)\) and \((\overline{4}, 1, 2)\) as the SM fermions while we consider the particles with quantum numbers \((4, 2, 1)\) and \((\overline{4}, 1, \overline{2})\) as exotic particles because these particles are distinguished by anomalous \( U(1) \) gauge symmetries (In the trinification models, these kinds of particles are obviously different.). With this convention, we can construct the three-family Pati-Salam models without tilted two-torus where the left-handed and right-handed SM fermions and the Higgs fields arise
from the intersections on the same two-torus, and then, we may explain three-family SM fermion masses and mixings.

In this kind of Pati-Salam model building, we require the intersection numbers to satisfy

$$I_{ab} = 3; \quad I_{ac} = -3; \quad I_{bc} \geq 1,$$

where $I_{ab} = 3$ and $I_{ac} = -3$ give us three families of the left-handed fermions and three families of the right-handed fermions, respectively, and $I_{bc} \geq 1$ gives us bidoublet Higgs field(s) with allowed Yukawa couplings.

We present one concrete model whose wrapping numbers and intersection numbers are given in Table IV. In this model, the absolute values of the intersection numbers on the second two-torus between the $U(4)_C$ and $U(2)_L$ stacks of D6-branes, between the $U(4)_C$ and $U(2)_R$ stacks of D6-branes, and between the $U(2)_L$ and $U(2)_R$ stacks of D6-branes are all three, and all the Yukawa couplings are allowed by anomalous $U(1)$ gauge symmetries. Therefore, we may explain the masses and mixings for three families of the SM fermions. Note that we have four additional D6-brane stacks $d$, $e$, $f$ and $g$ in this model, for which one can arbitrarily substitute them into filler brane stacks ($USp$ groups).

In general, the $U(4)_C \times U(2)_L \times U(2)_R$ gauge symmetry can be broken down to the $SU(3)_C \times SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$ gauge symmetry due to the G-S mechanism and the splittings of the $U(4)_C$ and $U(2)_R$ stacks of D6-branes. In our model, the $U(1)_{I_{3R}} \times U(1)_{B-L}$ gauge symmetry can only be broken down to the $U(1)_Y$ gauge symmetry by giving VEVs to the scalar components of the right-handed neutrino superfields or the neutral component in the multiplet $(\mathbf{4}, \mathbf{1}, \mathbf{2})$ from $I_{ac'}$ intersection. However, this Higgs mechanism can not preserve the D-flatness and F-flatness, and then breaks four-dimensional $N = 1$ supersymmetry. Therefore, the $U(1)_{I_{3R}} \times U(1)_{B-L}$ gauge symmetry breaking scale should be around the TeV scale.

C. $U(4)_C \times U(2)_L \times U(1)' \times U(1)''$ Model

In all the previous Pati-Salam like model building, the $U(1)_{I_{3R}}$ arises from the non-Abelian gauge symmetry, for example, $U(2)_R$ or $USp(2f)_R$. However, $U(1)_{I_{3R}}$ may come from a linear combination of $U(1)$ gauge symmetries.
TABLE IV: Wrapping numbers and intersection numbers in the $U(4)_C \times U(2)_L \times U(2)_R \times U(1)^4$ Model.

In our model building, we require

$$I_{ab} = 3 ; \quad I_{ac} = I_{ad} = -3 ; \quad I_{bc} \geq 1 ; \quad I_{bd} \geq 1 ,$$

(32)

where $I_{ab} = 3$ and $I_{ac} = I_{ad} = -3$ give us three families of the left-handed fermions and three families of the right-handed fermions, respectively, and $I_{bc} \geq 1$ gives us bidoublet Higgs fields with allowed Yukawa couplings.

Let us give a concrete supersymmetric model without the RR tadpoles and K-theory anomaly. We present the wrapping numbers and intersecting numbers in Table [V] and the spectrum in Tables [VIII] and [IX] in Appendix A. In particular, the c and d stacks of D6-branes are not T-dual to each other (If c and d stacks of D6-branes are T-dual to each other, the gauge symmetry in fact is $U(4)_C \times U(2)_L \times U(2)_R$). There are totally six $U(1)$ gauge symmetries where four combinations of them are global and their gauge fields obtain masses by the G-S mechanism. The rest two combinations, which are the massless anomaly-free $U(1)_{I_{3R}}$ and $U(1)_X$ gauge symmetries, are given by

$$U(1)_{I_{3R}} = \frac{1}{2}(U(1)_a + U(1)_b + 2U(1)_c) ,$$

(33)

$$U(1)_X = \frac{1}{2}(U(1)_a - U(1)_b + 2U(1)_d - 2U(1)_e - 2U(1)_f) .$$

(34)

In addition, the $U(4)_C \times U(2)_L \times U(1)' \times U(1)'' \times U(1)_e \times U(1)_f$ gauge symmetry can be broken down to the $SU(3)_C \times SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L} \times U(1)_X$ gauge symmetry by the G-S mechanism and the splitting of the $U(4)_C$ stack of D6-branes. Furthermore, the $U(1)_X$
gauge symmetry can be broken by giving VEVs to the SM singlets $1_d$ and $1_e$ (or $1_f$) which are charged under $U(1)_X$ (see the spectrum in Table IX in Appendix A). And the $U(1)_{I_{3R}} \times U(1)_{B-L}$ gauge symmetry can be broken down to the $U(1)_Y$ gauge symmetry by giving VEVs to $(\bar{4}_a, 1_e)$ (or $(\bar{4}_a, \bar{1}_e)$) and $(4_a, 1_f)$ (or $(4_a, \bar{1}_f)$). Because these Higgs mechanism can keep the D-flatness and F-flatness and then preserve four-dimensional $N = 1$ supersymmetry, these gauge symmetry breaking scales can be close to the string scale. However, only one family of the SM fermions can obtain masses.

| stack $N_a$ | $(n_1, l_1)(n_2, l_2)(n_3, l_3)$ | $A$ | $S$ | $b'$ | $c$ | $d'$ | $d$ | $e$ | $e'$ | $f$ | $f'$ | $f_{il}^{(1)}$ | $f_{il}^{(3)}$ |
|-------------|---------------------------------|-----|-----|------|-----|------|-----|-----|------|-----|------|----------------|----------------|
| $a$         | 8                               | -1  | 0   | 3    | 1   | -3   | -1  | -3  | -1   | 1   | 3    | -1            | 1              |
| $b$         | 4                               | (1,-1) | (1,2) | (1,0) | 2   | -2   | -2  | -2  | -2   | 2   | -2   | -2            | 2              |
| $c$         | 2                               | (1,1) | (1,0) | (1,-2)| -2  | 2    | -2  | -2  | -2   | 1   | 3    | 0             | 0              |
| $d$         | 2                               | (2,1) | (2,1) | (0,-1)| -6  | 6    | -6  | -6  | -6   | 1   | -1  | 3             | 3              |
| $e$         | 2                               | (-1,1) | (-1,0) | (1,-2)| -2  | 2    | -2  | -2  | -2   | -1  | 3    | 0             | -4             |
| $f$         | 2                               | (2,1) | (0,-1) | (2,1) | -6  | 6    | -6  | -6  | -6   | -1  | 3    | 0             | -4             |
| $f_{il}^{(3)}$ | 4                               | (0,1) | (0,0) | (0,-1)| -2  | 2    | -2  | -2  | -2   | -1  | 3    | 0             | -4             |
| $f_{il}^{(4)}$ | 4                               | (0,1) | (0,-1) | (1,0) | -2  | 2    | -2  | -2  | -2   | -1  | 3    | 0             | -4             |

TABLE V: Wrapping numbers and intersection numbers in the $U(4)C \times U(2)_L \times U(1)' \times U(1)'' \times U(1)_e \times U(1)_f \times USp(4) \times USp(4)$ model.

IV. TYPE IIB FLUX MODEL BUILDING

In this Section, we shall consider the trinification models, and the Pati-Salam like models on Type IIB orientifold with flux compactifications, which are very interesting because the supergravity fluxes can stabilize the dilaton and the complex structure parameters.

For the trinification models, we already have quite large RR charges due to the three $SU(3)$ groups. With Type IIB supergravity fluxes, the RR tadpole cancellation conditions are much more difficult to be satisfied. And in our detail calculations, we find that it may be impossible to find such a model.

For the Pati-Salam like models, the three-family and four-family Standard-like models with one unit of quantized flux and with the electroweak sector from $USp$ groups were
obtained \[29, 30\] by introducing magnetized D9-branes with large negative D3-brane charges in the hidden sector, and many supersymmetric and non-supersymmetric \(U(4)_C \times U(2)_L \times U(2)_R\) models were constructed by considering the magnetized D9-branes with large negative D3-brane charges in the SM observable sector \[32\]. Here, we consider a new flux model with \(U(4)_C \times U(2)_L \times U(2)_R\) gauge symmetry where the magnetized D9-branes with large negative D3-brane charges are introduced in the hidden sector. This kind of models has not been studied previously because it is very difficult to have supersymmetric D-brane configurations with more than three stacks of \(U(n)\) branes.

In the model building, we require the intersection numbers to satisfy the conditions in Eq. \(\text{[31]}\). We find a model with one unit of flux, and its wrapping numbers and intersection numbers are given in Table VI. Interestingly, no filler branes are needed so we do not have any \(USp\) groups. The two extra \(U(1)\) gauge symmetries are utilized to compensate the large positive D3-brane charges due to the supergravity fluxes. The \(U(4)_C \times U(2)_L \times U(2)_R\) gauge symmetry can be broken down to the \(SU(3)_C \times SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}\) gauge symmetry by the G-S mechanism and the splittings of the \(U(4)_C\) and \(U(2)_R\) stacks of D6-branes.

However, the \(U(1)_{I_{3R}} \times U(1)_{B-L}\) gauge symmetry can only be broken down to the \(U(1)_{Y}\) gauge symmetry at the TeV scale by giving VEVs to the scalar components of the right-handed neutrino superfields or the neutral component in the multiplet \((\overline{4}, 1, \overline{2})\) from \(I_{ac}\) intersection because this Higgs mechanism can not preserve the D-flatness and F-flatness, and then breaks four-dimensional \(N = 1\) supersymmetry. Also, we can only give masses to one family of the SM fermions.

| stack \(N_{a}(n_1, l_1)(n_2, l_2)(n_3, l_3)\) | \(A\) | \(b\) | \(b'\) | \(c\) | \(c'\) | \(d\) | \(d'\) | \(e\) | \(e'\) |
|---|---|---|---|---|---|---|---|---|---|
| \(a\) | 8 | (-1, 0) | (-1, 1) | ( 1, 1) | 0 | 0 | 3 | -3 | -1 | 3 | -3 | 3 | -3 |
| \(b\) | 4 | ( 1,-1) | ( 1, 2) | ( 1, 0) | 2 | -2 | - | - | 8 | 0(4) | 9 | -5 | 9 | -5 |
| \(c\) | 4 | ( 1, 1) | ( 1, 0) | ( 1,-2) | -2 | 2 | - | - | - | - | 5 | 5 | -9 | 5 | -9 |
| \(d\) | 2 | ( 2, 1) | (-2,-1) | ( 2, 1) | -54 | -10 | - | - | - | - | - | - | 0(0) | 64 |
| \(e\) | 2 | ( 2, 1) | (-2,-1) | ( 2, 1) | -54 | -10 | - | - | - | - | - | - | - | - |

**TABLE VI:** Wrapping numbers and intersection numbers in the \(U(4)_C \times U(2)_L \times U(2)_R \times U(1)^2\) model with one unit of flux.
V. DISCUSSIONS AND CONCLUSIONS

We have constructed new Standard-like models on Type II orientifolds. In Type IIA theory on $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold with intersecting D6-branes, we first constructed a three-family trinification model where all the SM fermions and Higgs fields belong to the bi-fundamental representations. The $U(3)_C \times U(3)_L \times U(3)_R$ gauge symmetry can be broken down to the $SU(3)_C \times SU(2)_L \times U(1)_{Y_L} \times U(1)_{I_{3R}} \times U(1)_{Y_R}$ gauge symmetry by the G-S mechanism and D6-brane splittings, and further down to the SM gauge symmetry at the TeV scale by the Higgs mechanism which can not preserve the D-flatness and F-flatness and then breaks four-dimensional $N = 1$ supersymmetry. In the general intersecting D-brane trinification models, the quark Yukawa couplings are allowed while the lepton Yukawa couplings are forbidden by anomalous $U(1)$ gauge symmetries. And in our model, we can only give masses to one family of the SM quarks. In addition, we constructed a Pati-Salam model which may generate suitable three-family SM fermion masses and mixings. The $U(4)_C \times U(2)_L \times U(2)_R$ gauge symmetry can be broken down to the $SU(3)_C \times SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$ gauge symmetry due to the G-S mechanism and D6-brane splittings. In our model, the $U(1)_{I_{3R}} \times U(1)_{B-L}$ gauge symmetry can only be broken down to the $U(1)_Y$ gauge symmetry at the TeV scale by Higgs mechanism. Moreover, we constructed for the first time a Pati-Salam like model with $U(4)_C \times U(2)_L \times U(1)' \times U(1)''$ gaug symmetry where the $U(1)_{I_{3R}}$ gauge symmetry comes from a linear combination of $U(1)$ gauge symmetries. And the $U(1)_{I_{3R}} \times U(1)_{B-L}$ gauge symmetry can be broken down to the $U(1)_Y$ gauge symmetry at (or close to) the string scale by Higgs mechanism. However, only one family of the SM fermions can obtain masses.

In Type IIB theory on $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold with flux compactifications, we could not construct the trinification model with supergravity fluxes because the three $SU(3)$ groups already contribute large RR charges. We constructed a new flux model with $U(4)_C \times U(2)_L \times U(2)_R$ gauge symmetry where the magnetized D9-branes with large negative D3-brane charges are introduced in the hidden sector. This kind of models has not been studied previously because it is very difficult to have supersymmetric D-brane configurations with more than three stacks of $U(n)$ branes. However, in our model, the $U(1)_{I_{3R}} \times U(1)_{B-L}$ gauge symmetry can only be broken down to the $U(1)_Y$ gauge symmetry at the TeV scale, and we can only give masses to one family of the SM fermions.
Acknowledgments

We would like to thank M. Cvetić for helpful discussions. T.L. is grateful to the George P. and Cynthia W. Mitchell Institute for Fundamental Physics for hospitality during the early stage of this project. The research of T.L. was supported by DOE grant DE-FG02-96ER40959, and the research of D.V.N. was supported by DOE grant DE-FG03-95-Er-40917.

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APPENDIX A: SPECTRA

In this Appendix, we present the spectrum in the $SU(3)_C \times SU(3)_L \times SU(3)_R \times U(1)^5 \times USp(2) \times USp(6)$ model with four global $U(1)$s from the G-S mechanism in Table VII and the spectrum in the $U(4)_C \times U(2)_L \times U(1)' \times U(1)'' \times U(1)_e \times U(1)_f \times USp(4) \times USp(4)$ model with anomaly free $U(1)_{I_{3R}}$ and $U(1)_X$ gauge symmetries in Tables VIII and IX.

| Rep. | Multi. | $U(1)_a$ | $U(1)_b$ | $U(1)_c$ | $U(1)_d$ | $U(1)_e$ | $U(1)_f$ | $USp(2)$ | $USp(6)$ |
|------|--------|----------|----------|----------|----------|----------|----------|----------|----------|
| $(3_a, \bar{3}_b)$ | 3 | 1 | -1 | 0 | 0 | 0 | -24 | 6 | 0 | 6 |
| $(\bar{3}_a, 3_c)$ | 3 | -1 | 0 | 1 | 0 | 0 | 12 | 6 | 0 | -12 |
| $(3_b, \bar{3}_e)$ | 3 | 0 | 1 | -1 | 0 | 0 | 12 | -12 | 0 | 6 |
| $(3_b, \bar{3}_c)$ | 1 | 0 | 1 | -1 | 0 | 0 | 12 | -12 | 0 | 6 |
| $(3_a, \bar{3}_b)$ | 1 | 1 | 1 | 0 | 0 | 0 | 0 | -6 | 0 | 6 |
| $(\bar{3}_a, 3_c)$ | 1 | -1 | 0 | -1 | 0 | 0 | 12 | -6 | 0 | 0 |
| $(\bar{3}_a, 1_d)$ | 3 | -1 | 0 | 0 | 1 | 0 | 10 | -4 | 2 | -10 |
| $(\bar{3}_a, 1_e)$ | 3 | -1 | 0 | 0 | 0 | 1 | 10 | -4 | 2 | -10 |
| $(\bar{3}_a, 1_e)$ | 3 | -1 | 0 | 0 | 0 | -1 | 14 | 4 | -2 | -2 |
| $(3_b, 1_d)$ | 6 | 0 | 1 | 0 | -1 | 0 | 14 | -2 | -2 | 4 |
| $(3_b, 1_e)$ | 6 | 0 | 1 | 0 | 0 | -1 | -2 | -2 | 4 | 4 |
| $(1_d, 1_e)$ | 16 | 0 | 0 | 0 | -1 | -1 | 4 | 8 | -4 | 8 |
| $A_a$ | 2 | -2 | 0 | 0 | 0 | 0 | 24 | 0 | 0 | -12 |
| $A_b$ | 2 | 0 | 2 | 0 | 0 | 0 | 24 | -12 | 0 | 0 |
| $S_a$ | 2 | 2 | 0 | 0 | 0 | 0 | -24 | 0 | 0 | 12 |
| $S_b$ | 2 | 0 | -2 | 0 | 0 | 0 | -24 | 12 | 0 | 0 |

Additional non-chiral and $USp(2) \times USp(6)$ Matter

TABLE VII: The spectrum in the $SU(3)_C \times SU(3)_L \times SU(3)_R \times U(1)^5 \times USp(2) \times USp(6)$ model with four global $U(1)$s from the G-S mechanism.
TABLE VIII: The SM fermions and Higgs fields in the $U(4)_{C} \times U(2)_{L} \times U(1)' \times U(1)'' \times U(1)_{c} \times U(1)_{f} \times USp(4) \times USp(4)$ model, with anomaly free $U(1)_{I_{3R}}$ and $U(1)_{X}$ gauge symmetries.
| Rep.     | Multi. | $U(1)_a$ | $U(1)_b$ | $U(1)_c$ | $U(1)_d$ | $U(1)_f$ | $2U(1)_{I_{3R}}$ | $2U(1)_X$ |
|----------|--------|----------|----------|----------|----------|----------|------------------|----------|
| $(4_a, 2_b)$ | 1 1 1 0 0 0 0 2 | 0 |
| $(4_a, 1_c)$ | 1 -1 0 -1 0 0 -3 -1 |
| $(4_a, 1_d)$ | 1 -1 0 0 -1 0 -1 -3 |
| $(4_a, 1_e)$ | 3 -1 0 0 0 1 -1 -3 |
| $(4_a, 1_f)$ | 1 -1 0 0 -1 0 -1 1 |
| $(4_a, 1_f)$ | 1 1 0 0 0 0 -1 1 3 |
| $(t_{2_b}, 1_d)$ | 3 1 0 0 0 0 1 1 -1 |
| $(t_{2_b}, 1_c)$ | 5 0 -1 0 -1 0 0 -1 -1 |
| $(t_{2_b}, 1_e)$ | 8 0 1 0 0 -1 0 1 1 |
| $(t_{2_b}, 1_f)$ | 3 0 -1 0 0 0 1 -1 -1 |
| $(t_{2_b}, 1_f)$ | 1 0 -1 0 0 0 -1 -1 3 |
| $(1_c, 1_d)$ | 1 0 0 1 -1 0 0 2 -2 |
| $(1_c, 1_d)$ | 3 0 0 1 1 0 0 2 2 |
| $(1_c, 1_f)$ | 5 0 0 1 0 0 -1 2 2 |
| $(1_c, 1_f)$ | 9 0 0 -1 0 0 -1 -2 2 |
| $(1_d, 1_c)$ | 1 0 0 0 -1 1 0 0 -4 |
| $(1_d, 1_c)$ | 3 0 0 0 1 1 0 0 0 |
| $(1_d, 1_f)$ | 16 0 0 0 -1 0 -1 0 0 |
| $(1_e, 1_f)$ | 5 0 0 0 0 1 -1 0 0 |
| $(1_e, 1_f)$ | 9 0 0 0 0 -1 -1 0 4 |
| $1_b$ | 2 0 2 0 0 0 0 2 -2 |
| $3_b$ | 2 0 -2 0 0 0 0 -2 2 |
| $1_c$ | 2 0 0 2 0 0 0 4 0 |
| $1_d$ | 6 0 0 0 2 0 0 0 4 |
| $1_e$ | 2 0 0 0 0 2 0 0 -4 |
| $1_f$ | 6 0 0 0 0 0 2 0 -4 |

**TABLE IX:** The extra particles in the $U(4)_C \times U(2)_L \times U(1)_I' \times U(1)' \times U(1)'' \times U(1)_c \times U(1)_f \times USp(4) \times USp(4)$ model, with anomaly free $U(1)_{I_{3R}}$ and $U(1)_X$ gauge symmetries.