# Standardlike models as type IIB flux vacua

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(Received 16 March 2005; published 26 May 2005)

We construct new semirealistic type IIB flux vacua on $Z_2 \times Z_2$ orientifolds with three- and four-standard model families and up to three units of quantized flux. The open string sector is comprised of magnetized D-branes and is T-dual to supersymmetric intersecting D6-brane constructions. The standard model sector contains magnetized D9-branes with negative D3-brane charge contribution. There are large classes of such models and we present explicit constructions for representative ones. In addition to models with one and two units of quantized flux, we also construct the first three- and four-family standard-like models with supersymmetric fluxes, i.e. comprising three units of quantized flux. Supergravity fluxes are due to the self-dual Neveu-Schwarz-Neveu-Schwarz and Ramond-Ramond three-form field strength and they fix the toroidal complex structure moduli and the dilaton. The supersymmetry conditions for the D-brane sector fix in some models all three toroidal Kähler moduli. We also provide examples where toroidal Kähler moduli are fixed by strong gauge dynamics on the ‘‘hidden sector’’ D7-brane. Most of the models possess Higgs doublet pairs with Yukawa couplings that can generate masses for quarks and leptons. The models have (mainly right-) chiral exotics.

DOI: 10.1103/PhysRevD.71.106008 PACS numbers: 11.25.Mj, 11.25.Wx

## I. INTRODUCTION

One of the challenging and essential problems in string theory is the construction of realistic string vacua, which can stabilize the moduli fields, generate standard model (SM)-like gauge structure and induce a (de Sitter) cosmological constant with supersymmetry breaking. Such constructions would provide a bridge between string theory and realistic particle physics. M-theory provides a framework where, in addition to perturbative heterotic string vacua, the physical string vacua could be probed in the perturbative type I, type IIA and type IIB superstring theory. In particular, the discovery of D-brane dynamics makes it possible to construct consistent four-dimensional supersymmetric $N = 1$ chiral models with non-Abelian gauge symmetry on type II orientifolds, by employing conformal field theory techniques in the open string sector. The first such supersymmetric models were based on $Z_2 \times Z_2$ orientifolds [1,2]. [Nonsupersymmetric constructions were given in [3–6] (see also [7] and for earlier work [8,9]).] Subsequently, a number of SM-like models, grand unified theory (GUT) models, and their variations have been constructed in various orbifold backgrounds, and the associated phenomenology has been discussed. (For a partial list of nonsupersymmetric constructions, see [10], further supersymmetric constructions are given in [11–14] and further developments in connection with the study of effective couplings and phenomenological implications, see [15–21] and references therein.)

(Recently important progress has been made in constructions of supersymmetric chiral solutions of type II Gepner models; see [22,23] and references therein. Specifically, the recent impressive results of [23] provide large classes of three-family standard-like models with no chiral exotics. Note however, that these exact conformal field theory models are located at the special points in the moduli space where the geometric picture is lost. In particular couplings, such as Yukawa couplings, do not possess hierarchies associated with the size of the internal spaces, such as in the case of the toroidal orbifolds with D-branes. In addition, due to the lack of geometric interpretation, the introduction of supergravity fluxes does not seem to be possible.)

In spite of these successes, the moduli stabilization in open string and closed string sectors remained an open problem, even though in some cases some complex structure moduli (in the type IIA picture) and dilaton fields may be stabilized due to nonperturbative gauge dynamics, associated with the gaugino condensation in the hidden sector (see, e.g., [24].) Turning on supergravity fluxes introduces a supergravity potential, which provides another way to stabilize the compactification moduli fields by lifting continuous moduli space of the string vacua in the effective four-dimensional theory (see, e.g., [25]). However, the introduction of supergravity fluxes imposes strong constraints on consistent constructions, since such fluxes modify the global Ramond-Ramond (RR) tadpole cancellation conditions. Meanwhile, the fluxes will typically generate a back-reaction on the original geometry of the internal space, thus changing the nature of the internal space.

On the type IIA side the supersymmetry conditions of flux compactifications are less understood. Nevertheless recent work [26,27] revealed the existence of unique flux vacua for massive type IIA string theory with SU(3) structure, whose geometry of the internal six-dimensional space is nearly Kähler and four-dimensional space is anti de Sitter (AdS) (for the discussion on necessary and sufficient conditions of $N = 1$ compactifications of massive IIA supergravity to AdS(4) with SU(3) structure, see also [28]). One such example is the $\frac{\mathrm{SU}(2)^3}{\mathrm{SU}(2)} \simeq S_3 \times S_3$ coset space that has three supersymmetric three-cycles that add up to zero in homology [27,29]. Therefore the total charge of the D6-branes wrapping such cycles is zero and no introduction of orientifold planes on such spaces is needed. Moreover, since the three-cycles intersect pairwise, the massless chiral matters appear at these intersections. This construction [27] therefore provides an explicit example of supersymmetric flux compactifications with intersecting D6-branes. Further progress also has been made in the construction of $N = 1$ supersymmetric type IIA flux vacua with SU(2) structures [27], leading to examples with the internal space conformally Calabi-Yau. However, explicit constructions of models with intersecting probe D6-branes for such flux compactifications is still awaiting further study.

On the type IIB side the intersecting D6-brane constructions correspond to models with magnetized branes with the role of the intersecting angles played by the magnetic fluxes on the branes. The dictionary for the consistency and supersymmetry conditions between the two T-dual constructions is straightforward, see e.g., [30,31]. The supersymmetric type IIB flux compactifications are also better understood; see, e.g., [32–35], [30,31], and references therein. In particular, examples of supersymmetric fluxes and the internal space conformally Calabi-Yau are well known. The prototype example is a self-dual combination of the Neveu-Schwarz-Neveu-Schwarz (NS-NS) $H_3$ and RR $F_3$ three-forms, corresponding to the primitive (2,1) form on Calabi-Yau space. Since the backreaction of such flux configurations is mild, i.e., the internal space remains conformal to Calabi-Yau, these type IIB flux compactifications are especially suitable for adding the probe magnetized D-branes in this background. However, the quantization conditions on fluxes and the modified tadpole cancellation conditions constrain the possible D-brane configurations severely. In Refs. [30,31] the techniques for consistent chiral flux compactifications on orbifolds were developed, however, no explicit supersymmetric chiral SM-like models were obtained.

Most recently, by introducing magnetized D9-branes carrying negative D3-brane charges in the hidden sector, in Ref. [36] the first example of three-family SM-like string vacuum with one unit of quantized flux turned on was obtained, and subsequently, the first four-family SM-like string vacuum with one unit of fluxes was constructed in [37]. These constructions could be T-dual to the supersymmetric models of intersecting D6-branes on $Z_2 \times Z_2$ orientifold with the $\mathrm{Sp}(2)_L \times \mathrm{Sp}(2)_R$ or $\mathrm{Sp}(2f)_L \times \mathrm{Sp}(2f)_R$ gauge symmetry in the electroweak sector, respectively [14,18]. [Without fluxes, the first models of that type were toroidal models with intersecting D6-branes [18] where the RR tadpoles were not explicitly cancelled. (For the subsequent generalization to tilted tori see [38].) The $Z_2 \times Z_2$ orientifold construction in [14] provided the first model of that type that cancelled RR tadpoles by introducing an additional stack of D6-branes with unitary symmetry; in the T-dual picture those are precisely the magnetized D9-branes with the negative contribution to the D3-brane charge.] In spite of these successes, we are confronted by a number of problems:

(i) The semirealistic SM-like string vacua with four-dimensional $N = 1$ supersymmetric fluxes have not been constructed, yet. In view of this drawback, effects of nonsupersymmetric fluxes, as a key mechanism for breaking supersymmetry, has been addressed [39–42]. This analysis [39] leads to soft supersymmetry breaking masses $M_{\mathrm{soft}} \sim \frac{M_s^2}{M_{\mathrm{Pl}}}$ where $M_s$ and $M_{\mathrm{Pl}}$ are the string scale and Planck scale, respectively. In the toroidal orbifold constructions of this type $M_s \sim M_{\mathrm{Pl}}$. In order to achieve $M_{\mathrm{soft}} \ll M_{\mathrm{Pl}}$ and the stabilization of the electroweak scale, one has to introduce inhomogeneous warp factors in the internal space, which is hard to realize for toroidal orbifold compactifications.

(ii) These flux vacua stabilize the dilaton and toroidal complex structure moduli. However, the Kähler moduli do not enter the flux-induced superpotential and hence are hard to be completely fixed. So, we are still typically faced with the vacuum degeneracy problem. [The Kähler-moduli fields in type IIB string theory are T-dual to the complex structure moduli in type IIA string theory (intersecting D6-brane scenario). In the T-dual picture, the latter moduli can often be stabilized by employing nonperturbative dynamical mechanism, such as the gaugino and matter condensation in the hidden sector. However, these mechanisms are difficult to employ on the type IIB side due to the additional matter content on the associated magnetized D-branes.]

(iii) For explicit type IIB orientifolds the imaginary self-dual fluxes are quantized in rather large flux units, e.g., for $Z_2 \times Z_2$ the orientifolds elementary flux unit is 64. Therefore the constructions of semirealistic flux vacua is very constrained; only one unit of flux is allowed for known semirealistic three- [36] and four-family [37] models. Thus the introduction of fluxes is restricted to very few known semirealistic examples and is not typical.

Following our previous work [37], in this paper we systematically study new constructions of three- and four-family SM-like string vacua with supergravity fluxes on type IIB $Z_2 \times Z_2$ orientifolds. The major technical difficulty in constructions of semirealistic flux vacua on type IIB orientifolds is to ensure the cancellation of the large positive D3-brane charge contribution to RR tadpoles by the fluxes. Similar to the D-brane models without fluxes [14], and the subsequent work with fluxes [36,37], the important role in tadpole cancellation is played (in the type IIB picture) by magnetized D9-branes which carry large negative D3-brane charges. In the past constructions ([14,36,37]) such D9-branes were introduced as a part of the "hidden sector."

In this paper, we consider new types of constructions where *the magnetized D9-branes with large negative D3-brane charges are introduced as a part of the SM sector.* For this new setup, we find that the constructions of SM-like flux vacua are much less constrained and obtained a large class of new models. In particular, in addition to many new models with one unit of quantized flux, *we obtain first three- and four-family models with two units of quantized charge, as well as the first three- and four-family examples with supersymmetric flux, i.e. three units of quantized flux*. Such supersymmetric three- and four-family SM-like models have toroidal Kähler moduli fixed by supersymmetry conditions [43] and the string scale can be close to the Planck scale. These models have (mainly right-) chiral exotics. However, most of the models have Higgs doublet pairs with Yukawa couplings to quarks and leptons and thus can generate the one-family SM fermion masses, or the suitable fermion masses and mixings for two families, and even give large masses to some of the bi-fundamental chiral exotics at the tree level. The inclusion of quantum corrections for the latter case may allow for fermionic mass spectra with hierarchy and mixings for all three families. Finally, with the open string moduli fixed by the flux-induced soft masses, we are able to construct the first SM-like string vacua with strong infrared gauge dynamics on the "hidden sector" D7-branes, which can generate the Veneziano-Yankielowicz superpotential, and hence may help stabilize the toroidal Kähler moduli à la KKLT mechanism [44].

The paper is organized in the following way: In Sec. II, we systematize the constructions of supersymmetric string vacua with supergravity fluxes on type IIB $Z_2 \times Z_2$ orientifolds. In Sec. III, we classify the classes of the SM-like flux vacua on type IIB $Z_2 \times Z_2$ orientifolds. Subsequently, we discuss in detail explicit constructions of the representative models with one, two, and three (supersymmetric) units of quantized flux, as well as a specific construction with gaugino condensation on the "hidden sector" D7-branes. In the appendix we provide tables of all the explicitly constructed representative models. We conclude with discussions and open problems in Sec. IV.

## II. MAGNETIZED D-BRANES AND TYPE IIB FLUX COMPACTIFICATIONS ON $T^6/(Z_2 \times Z_2)$ ORIENTIFOLDS

Flux compactifications on simplest toroidal $T^6$ type IIB orientifolds, on which most of the previous work has focused, are unlikely to provide a framework for constructions of semirealistic flux vacua. We shall therefore focus on the simplest orbifold constructions, i.e. on $T^6/(Z_2 \times Z_2)$ orientifolds.

The internal space $T^6$ is chosen to be factorized as a direct product of three two-tori, i.e. $T^6 = T^2 \times T^2 \times T^2$, whose complex coordinates are $z_i$, $i = 1, 2, 3$ for the $i$th two-torus, respectively. The generators $\theta$ and $\omega$ for the orbifold group $Z_2 \times Z_2$, act on the complex coordinates of $T^6$ as

$$\begin{aligned}\theta&:(z_1, z_2, z_3) \rightarrow (-z_1, -z_2, z_3),\\ \omega&:(z_1, z_2, z_3) \rightarrow (z_1, -z_2, -z_3).\end{aligned} \tag{1}$$

On $Z_2 \times Z_2$ orbifold type IIB string theory contains in the untwisted sector the four-dimensional $N = 2$ supergravity multiplet, the dilaton hypermultiplet, $h_{11}$ hypermultiplets, and $h_{21}$ vector multiplets, which are all massless. The orbifold action projects out several components of the metric of a general $T^6$ geometry and, as a result, we are left with fewer Kähler and complex structure parameters. These are encoded for the untwisted moduli in terms of the Hodge numbers, as $(h_{11}, h_{21})_{\mathrm{unt}} = (3, 3)$. On the other hand, each of the three elements $\theta$, $\omega$, and $\theta\omega$ has a fixed-point set given by 16 $T^2$'s, and the corresponding twisted sectors also contribute to the Hodge numbers of the orbifold. For a particular choice of discrete torsion, this contribution is given by $(h_{11}, h_{21})_{\mathrm{tw}} = (0, 3 \times 16)$. The contributions from both the untwisted and twisted sectors hence add up to $(h_{11}, h_{21}) = (3, 51)$.

Orientifold planes are necessary for the introduction of the open string sector, and the associated orientifold projection can be denoted by $\Omega R$, where $\Omega$ is the world-sheet parity projection and $R$ (acting on type IIA as the holomorphic $Z_2$ involution) acts on the complex coordinates as

$$R:(z_1, z_2, z_3) \rightarrow (-z_1, -z_2, -z_3). \tag{2}$$

Thus, the model contains 64 $O3$-planes and 4 $O7_i$-planes, transverse to the $i$th two-torus. This orientifold action projects the above $N = 2$ spectrum to an $N = 1$ supergravity multiplet, the dilaton chiral multiplet, and 6 untwisted and 48 twisted geometrical chiral multiplets.

In order to cancel the negative RR charge contributions, due to these $O$-planes, we need to introduce D$(3 + 2n)$-branes which are filling up four-dimensional Minkowski space-time and wrapping $2n$-cycles on the compact manifold. We choose the construction with magnetized D-branes. (A detailed discussion for toroidal/orbifold compactifications with magnetized D-branes is given, e.g., in [30].) Concretely, for one stack of $N_a$ D-branes wrapped $m_a^i$ times on the $i$th two-torus $T_i^2$, we turn on $n_a^i$ units of magnetic fluxes $F_a$ for the center of mass $\mathrm{U}(1)_a$ gauge factor on each $T_i^2$, such that

$$m_a^i \frac{1}{2\pi} \int_{T_i^2} F_a^i = n_a^i. \tag{3}$$

Hence, the topological information of this stack of D-branes is encoded in $N_a$-number of D-branes and the co-prime number pairs $(n_a^i, m_a^i)$. The D9-, D7-, D5-, and D3-branes contain 0, 1, 2, and 3 vanishing $m_a^i$s, respectively. Introducing for the $i$th two-torus the even homology classes $[\mathbf{0}_i]$ and $[\mathbf{T}_i^2]$ for the point and the two-torus, respectively, the vectors of RR charges of the $a$th stack of D-branes and its image are

$$[\Pi_a] = \prod_{i=1}^{3}(n_a^i[\mathbf{0}_i] + m_a^i[\mathbf{T}_i^2]),$$
$$[\Pi_a'] = \prod_{i=1}^{3}(n_a^i[\mathbf{0}_i] - m_a^i[\mathbf{T}_i^2]), \tag{4}$$

respectively. Similarly, for the $O3$- and $O7_i$-planes appearing on $T^6/(Z_2 \times Z_2)$ orientifold which, respectively, correspond to $\Omega R$, $\Omega R\omega$, $\Omega R\theta\omega$, and $\Omega R\theta$ $O$-planes, we have

$$\begin{aligned} \Omega R&:[\Pi_{O3}] = [\mathbf{0}_1] \times [\mathbf{0}_2] \times [\mathbf{0}_3]; \\ \Omega R\omega&:[\Pi_{O7_1}] = -[\mathbf{0}_1] \times [T_2^2] \times [T_3^2]; \\ \Omega R\theta\omega&:[\Pi_{O7_2}] = -[T_1^2] \times [\mathbf{0}_2] \times [T_3^2]; \\ \Omega R\theta&:[\Pi_{O7_3}] = -[T_1^2] \times [T_2^2] \times [\mathbf{0}_3]. \end{aligned} \tag{5}$$

The "intersection numbers," which determine the chiral massless spectrum, are

$$\begin{aligned} I_{ab} &= [\Pi_a] \cdot [\Pi_b] = \prod_{i=1}^{3}(n_a^i m_b^i - n_b^i m_a^i), \\ I_{ab'} &= [\Pi_a] \cdot [\Pi_{b'}] = -\prod_{i=1}^{3}(n_a^i m_b^i + n_b^i m_a^i), \\ I_{aa'} &= [\Pi_a] \cdot [\Pi_{a'}] = -8\prod_{i=1}^{3}(n_a^i m_a^i), \\ I_{aO} &= \sum_p [\Pi_a] \cdot [\Pi_{Op}] \\ &= 8(-m_a^1 m_a^2 m_a^3 + m_a^1 n_a^2 n_a^3 + n_a^1 m_a^2 n_a^3 + n_a^1 n_a^2 m_a^3), \end{aligned} \tag{6}$$

where $[\Pi_{Op}] = [\Pi_{O3}] + [\Pi_{O7_1}] + [\Pi_{O7_2}] + [\Pi_{O7_3}]$ is the sum of $O3$-plane and $O7_i$-plane homology classes.

TABLE I. General spectrum for magnetized D-branes on the type IIB $T^6/(Z_2 \times Z_2)$ orientifold. The representations in the table refer to $U(N_a/2)$, the resulting gauge symmetry due to $Z_2 \times Z_2$ orbifold projection. For supersymmetric constructions, scalars combine with fermions to form chiral supermultiplets.

| Sector | Representation |
|---|---|
| $aa$ | $U(N_a/2)$ vector multiplet |
| | 3 adjoint chiral multiplets |
| $ab + ba$ | $I_{ab}$ $(\square_a, \overline{\square}_b)$ fermions |
| $ab' + b'a$ | $I_{ab'}$ $(\square_a, \square_b)$ fermions |
| $aa' + a'a$ | $\frac{1}{2}(I_{aa'} - \frac{1}{2}I_{a,Op})$ $\square\square$ fermions |
| | $\frac{1}{2}(I_{aa'} + \frac{1}{2}I_{a,Op})$ $\begin{smallmatrix}\square\\\square\end{smallmatrix}$ fermions |

Similar to the discussions in [2], the physical chiral spectrum should be invariant under the full orientifold symmetry group and is tabulated in Table I. Flux vacua on type IIB orientifolds with four-dimensional $N = 1$ supersymmetry are primarily constrained by the RR tadpole cancellation conditions and conditions for $N = 1$ supersymmetry in four dimension, which we describe in detail in the following subsections.

### A. RR tadpole cancellation conditions

In the type IIB picture the fluxes, we consider, are due to the self-dual three-from field strength, which contributes to the D3-brane field strength equation of motion, and thus modifies the D3 charge conservation (RR-tadpole cancellation) conditions on the compact orientifold. The RR charges, carried by magnetized D-branes, are classified by their associated homology classes. Explicitly, for one stack of $N_a$ D-branes with wrapping numbers $(n_a^i, m_a^i)$, it carries D3-, D5-, D7-, and D9-brane RR charges

$$\begin{aligned} Q3_a &= N_a n_a^1 n_a^2 n_a^3, \qquad (Q5_i)_a = N_a m_a^i n_a^j n_a^k, \\ (Q7_i)_a &= N_a n_a^i m_a^j m_a^k, \qquad Q9_a = N_a m_a^1 m_a^2 m_a^3, \end{aligned} \tag{7}$$

where $i \neq j \neq k$, and a permutation is implied for $(Q5_i)_a$ and $(Q7_i)_a$. So, the RR tadpole cancellation conditions can be described as

$$\sum_a N_a[\Pi_a] + \sum_a N_a[\Pi_{a'}] + \sum_p N_{Op} Q_{Op}[\Pi_{Op}] + N_{\rm flux} = 0, \tag{8}$$

where the third term contribution comes from the $O3$- and $O7_i$-planes, with $N_{Op}$ and $Q_{Op}$ denoting their numbers and RR charges, respectively. And $N_{\rm flux}$ is the amount of the fluxes turned on and is quantized in units of the elementary flux.

For a supersymmetric Dp/Dp′-brane system on type IIB $T^6/(Z_2 \times Z_2)$ orientifold, only D3- and D7-branes are allowed to be wrapped along the orientifold planes. (In the following we shall refer to this type of branes as "filler branes.") Given that $N_{Op}Q_{Op} = 2^{9-p}(-2^{p-4}) \equiv -32$ in Dp-brane units for Sp-type $O$-planes, the RR tadpole cancellation conditions can be further simplified as

$$\begin{aligned} -N^{(0)} - \sum_a Q3_a - \frac{1}{2}N_{\rm flux} &= -16, \\ -N^{(i)} + \sum_a (Q7_i)_a &= -16, \; i \neq j \neq k, \end{aligned} \tag{9}$$

where $N^{(0)}$ and $N^{(i)}$ with $i = 1$, 2, and 3, respectively, denote the number of filler branes, i.e. D-branes which wrap along the $O3$- and $O7_i$-planes and only contribute to one of the four kinds of D3- and D7-brane charges. As for D5- and D9-brane RR tadpoles, their cancellations are automatic since these D-branes and their $\Omega R$ images carry the same absolute value of the corresponding charges but with opposite sign of charges.

### B. Conditions for four-dimensional $N = 1$ supersymmetry

Four-dimensional $N = 1$ supersymmetric vacua from flux compactification require that $1/4$ supercharges of the ten-dimensional (T-dual) type I theory be preserved in both open string and close string sectors. We shall discuss both sectors separately.

For the closed string sector, the specific type IIB flux solution on orientifolds comprises of self-dual three-form field strength and it has been discussed, e.g., in [32,45]. While RR $F_3$ and NS-NS $H_3$ three-form fluxes are turned on, the induced three-form $G_3 = F_3 - \tau H_3$, with $\tau = a + i/g_s$ being type IIB axion-dilaton coupling, contributes to the D3-brane RR charges

$$N_{\mathrm{flux}} = \frac{1}{(4\pi^2\alpha')^2}\int_{X_6} H_3 \wedge F_3 = \frac{1}{(4\pi^2\alpha')^2}\frac{i}{2\tau_I}\int_{X_6} G_3 \wedge \bar{G}_3, \tag{10}$$

where $\tau_I$ is the imaginary part of the complex coupling $\tau$. Dirac quantization conditions of $F_3$ and $H_3$ on $T_6/(Z_2 \times Z_2)$ orientifold require that $N_{\mathrm{flux}}$ be a multiple of 64, and the Bogomolnyi-Prasad-Sommerfield-like self-duality condition: $*_6 G_3 = iG_3$ ensures that its contribution to the RR charges is positive. Supersymmetric configuration implies that $G_3$ background field should be a primitive self-dual (2,1) form. A specific supersymmetric solution which is useful for our purpose is [30,45]

$$G_3 = \frac{8}{\sqrt{3}}e^{-\pi i/6}(d\bar{z}_1 dz_2 dz_3 + dz_1 d\bar{z}_2 dz_3 + dz_1 dz_2 d\bar{z}_3), \tag{11}$$

where the additional factor 4 is due to the $Z_2 \times Z_2$ orbifold symmetry. The fluxes stabilize the complex structure toroidal moduli at values

$$\tau_1 = \tau_2 = \tau_3 = \tau = e^{2\pi i/3}, \tag{12}$$

leading to the RR tadpole contribution in Eq. (9)

$$N_{\mathrm{flux}} = 192. \tag{13}$$

This result therefore implies that in order to construct supersymmetric SM-like flux vacua, we have to introduce at least the D3 charge conservation condition, which is thus hard to achieve.

In the open string sector, for D-branes with world-volume magnetic field $F^i = \frac{n^i}{m^i\chi^i}$, the four-dimensional $N = 1$ supersymmetry is ensured if and only if $\sum_i \theta_i = 0$ (mod $2\pi$) is satisfied [30]. Here the "angle" $\theta_i$ with the range $\{0, 2\pi\}$ is determined in terms of the world-volume magnetic field as $\tan(\theta_i) \equiv (F^i)^{-1} = \frac{m^i\chi^i}{n^i}$ and $\chi^i = R_1^i R_2^i$; the area of the $i$th two-torus $T_i^2$ in $\alpha'$ units is the Kähler modulus of the $i$th two-torus $T_i^2$. This supersymmetry condition can be cast in the form: $\sum_i (F^i)^{-1} - (F^1F^2F^3)^{-1} = 0$, along with $\sum_{i<j}(F^iF^j)^{-1} - 1 < 0$ for $n_in_jn_k > 0$ or $\sum_{i<j}(F^iF^j)^{-1} - 1 > 0$ for $n_in_jn_k < 0$, which can be rewritten in the following form:

$$-x_A Q9_a + x_B(Q5_1)_a + x_C(Q5_2)_a + x_D(Q5_3)_a = 0,$$
$$-Q3_a/x_A + (Q7_1)_a/x_B + (Q7_2)_a/x_C + (Q7_3)_a/x_D < 0, \tag{14}$$

where $x_A = \lambda$, $x_B = \lambda/\chi^2\chi^3$, $x_C = \lambda/\chi^1\chi^3$, $x_D = \lambda/\chi^1\chi^2$. The positive parameter $\lambda$ has been introduced to put all the variables $Q9$, $Q7_i$, $Q5_i$, and $Q3_i$ on equal footing. These supersymmetry conditions can be cast easily in the T-dual form of the type IIA supersymmetry constraints discussed in [2].

## III. CONSTRUCTIONS OF SM-LIKE STRING VACUA FROM TYPE IIB FLUX COMPACTIFICATION

Similar to the past constructions (see specifically those of [13]), we construct the SM-like models as descendants of the Pati-Salam model based on $\mathrm{SU}(4)_C \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$ gauge symmetry. The hypercharge is

$$Q_Y = Q_{I_{3R}} + \frac{Q_{B-L}}{2}, \tag{15}$$

where the nonanomalous $\mathrm{U}(1)_{B-L}$ is obtained from the splitting of the $\mathrm{U}(4)_C$ branes, i.e. $\mathrm{U}(4) \to \mathrm{U}(3)_C \times \mathrm{U}(1)_{B-L}$. Similarly, the anomaly-free $\mathrm{U}(1)_{I_{3R}}$ gauge symmetry is from the non-Abelian $\mathrm{U}(2)_R$ or $\mathrm{Sp}(2)_R$ gauge symmetry. There are three main frameworks to realize the Pati-Salam gauge sector in the type IIB magnetized D-brane scenario (T-dual to the type IIA intersecting D6-brane one):

(i) The starting observable sector gauge symmetry is $\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{U}(2)_R$ [13]. In this framework, the three "anomalous" gauge symmetries $\mathrm{U}(1)_C$, $\mathrm{U}(1)_L$, and $\mathrm{U}(1)_R$ can be treated as global ones, since the associated gauge bosons obtain masses via $B \wedge F$ Chern-Simons couplings. [Those are effective couplings that arise from the D-brane world-volume Chern-Simons couplings and are responsible for the Abelian gauge anomaly cancellation via the Green-Schwarz mechanism.] The gauge symmetry breaking chain is of the form:

$$\begin{aligned}&\mathrm{SU}(4) \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R\\ \to&\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L}\\ \to&\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_{I_{3R}} \times \mathrm{U}(1)_{B-L}\\ \to&\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y,\end{aligned} \tag{16}$$

where the first and second step can be achieved by splitting the $U(4)_C$ branes and $U(2)_R$ branes in one two-torus direction and the third step by giving vacuum expectation values (VEVs) to the scalar components of right-handed neutrino chiral superfields (or a scalar component of an exotic nonchiral chiral superfields) at the TeV scale. Alternatively, we may skip the second step and directly break the $SU(2)_R \times U(1)_{B-L}$ gauge symmetry down to the $U(1)_Y$ by giving VEVs to the scalar components of right-handed neutrino chiral superfields. Within this framework one typically obtains enough SM Higgs doublet pairs (from the chiral or the nonchiral massless sector) with Yukawa couplings to quarks and leptons and hence may generate the fermion masses and mixings in the SM-sector at the tree level. However, it also should be noted that in a few cases all the SM Higgs doublet pairs have the global U(1) quantum numbers that do not allow for Yukawa couplings to quarks and leptons; these models therefore face serious phenomenological obstacles.

(ii) The starting gauge symmetry is one-family $U(4) \times Sp(8)_L \times Sp(8)_R$ or two-family $U(4) \times Sp(4)_L \times Sp(4)_R$, which can be broken down to the four-family $U(4) \times U(2)_L \times U(2)_R$ or $U(4) \times SU(2)_L \times SU(2)_R$ by parallel splitting the D-branes, originally positioned on the $O$-planes, in three or two-tori directions, respectively. [Both the string theory and field theory aspects of the brane splittings in this framework are discussed in detail in [14] for constructions without fluxes. The flux vacua of that type (with one unit of quantized flux) was constructed in [37].] In the field theory picture ("Higgsing") the four-families ($f = 4$) are obtained when we decompose the original chiral supermultiplets $(4, 8, 1)$ and $(\bar{4}, 1, 8)$, or $(4, 4, 1)$ and $(\bar{4}, 1, 4)$ into four copies of $(4, 2, 1)$ and $(\bar{4}, 1, 2)$ after the gauge symmetry breaking. These Higgsings, as discussed in [14], preserve the D- and F-flatness, and thus the symmetry breaking can take place at the string scale. Thus, in these cases the resulting spectrum is that of four-dimensional $N = 1$ supersymmetric four-family Pati-Salam models. The symmetry breaking chains for these two pictures are given, respectively, by

$$\begin{aligned} & SU(4) \times Sp(8)_L \times Sp(8)_R \\ \rightarrow & SU(4) \times U(2)_L \times U(2)_R \\ \rightarrow & SU(3)_C \times U(2)_L \times U(2)_R \times U(1)_{B-L} \\ \rightarrow & SU(3)_C \times U(2)_L \times U(1)_Y. \end{aligned} \tag{17}$$

and

$$\begin{aligned} & SU(4) \times Sp(4)_L \times Sp(4)_R \\ \rightarrow & SU(4) \times SU(2)_L \times SU(2)_R \\ \rightarrow & SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \\ \rightarrow & SU(3)_C \times SU(2)_L \times U(1)_Y. \end{aligned} \tag{18}$$

The first and second step can be achieved by splitting the Sp- and U(4)-branes at the string scale, and the third step by giving VEVs to the scalar components of the right-handed neutrino superfields at the TeV scale. Note that for the model with the original $U(4) \times Sp(8)_L \times Sp(8)_R$ symmetry, the resulting $U(1)_L$ and $U(1)_R$ are not anomalous since they are part of the non-Abelian Sp symmetries. One expects that at least the gauge boson of $U(1)_L$ will obtain a mass at the electroweak scale, which is excluded by experiments. In order to evade this problem, we allow only for the $Sp(8)_R$ gauge symmetry in the SM sector, while the $Sp(8)_L$ gauge symmetry is not. Moreover, we consider the variants of the $U(4) \times Sp(4)_L \times Sp(4)_R$ model, i.e. the models with the gauge symmetry $U(4) \times U(2)_L \times Sp(4)_R$ or $U(4) \times Sp(4)_L \times U(2)_R$. [This analysis also can be applied to three-family models with gauge symmetry $U(4) \times Sp(6)_L \times Sp(6)_R$ where the $Sp(6)_L \times Sp(6)_R$ gauge symmetry can be broken down to the $Sp(2)_L \times Sp(2)_R$ by the Higgs mechanism, however this symmetry breaking pattern breaks supersymmetry [14] and thus may only be implemented within the framework of supersymmetry breaking at scale larger than the electroweak scale.]

(iii) The starting symmetry is the Pati-Salam-like $U(4)_C \times Sp(2)_L \times Sp(2)_R$. Without fluxes, the first models of that type were toroidal orientifolds with intersecting D6-branes [18] where the RR tadpoles were not explicitly cancelled. The $Z_2 \times Z_2$ orientifold construction in Ref. [14] cancelled the RR-tadpoles by introducing an additional stack of branes with unitary symmetry in the hidden sector. In the T-dual type IIB picture those are the magnetized D9-branes with *a negative* D3 *charge*. In Ref. [36] these types of magnetized D9-branes were employed to find the MSSM-like model with one unit of quantized flux turned on. In this framework, the starting gauge symmetry is similar to the framework (i); the initial framework gauge symmetry $U(4)_C \times Sp(2)_L \times Sp(2)_R$ can be broken down to the SM gauge symmetry as

$$\begin{aligned} & SU(4) \times SU(2)_L \times SU(2)_R \\ \rightarrow & SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \\ \rightarrow & SU(3)_C \times SU(2)_L \times U(1)_Y, \end{aligned} \tag{19}$$

where the first step can be achieved again by splitting the U(4) D-branes at string scale, and the second step is achieved by giving VEVs to the scalar components of the right-handed neutrino chiral superfields at the TeV scale. Note that in the present case there is no $U(1)_L$. However, in the models constructed in Ref. [18,36], there exists only one SM Higgs doublet pair; in this case a generic problem is that only the third family can obtain the tree-level masses and it is difficult to give masses to the first two families at the quantum level [18]. In addition, the $SU(2)_R \times U(1)_{B-L}$ gauge symmetry can be broken down to the $U(1)_Y$ only by giving VEVs to the scalar components of the right-handed neutrino chiral superfields at the TeV scale in this kind of model.

The Pati-Salam-type models with only one SM Higgs pair suffer from serious phenomenological problems [46]. Even though one may be able to generate the most general Yukawa couplings via radiative corrections, the mass matrix of up-type quarks is proportional to that of down-type quarks, and the mass matrix of neutrinos is proportional to that of leptons. (Note the renormalization group equation running effects on these mass matrices are negligible.) Therefore, on the one hand, the masses for the quarks, leptons, and neutrinos satisfy

$$m_u:m_c:m_t = m_d:m_s:m_d, \qquad m_{\nu_e}:m_{\nu_\mu}:m_{\nu_\tau} = m_e:m_\mu:m_\tau. \tag{20}$$

The above fermion mass relations are obviously wrong from the known experiments. On the other hand, the Cabibbo-Kobayashi-Maskawa quark-mixing matrix and the Pontecorvo-Maki-Nakagawa-Sakata neutrino mixing matrix are proportional to the identity matrix which implies that the quark and neutrino mixing angles vanish, again in contradiction with experiments. Note that these problems for the fermion masses and mixings cannot be solved by loop corrections because the $SU(2)_R \times U(1)_{B-L}$ gauge symmetry is broken at TeV scale. There is a consensus that the minimal supersymmetric Pati-Salam or left-right model should have at least two SM Higgs doublets [48]; therefore, the construction (without fluxes) in [14], which actually contains two SM Higgs doublets, is the construction which can realize the embedding of the supersymmetric Pati-Salam model with realistic features in the type IIA intersecting D-brane scenario, or equivalently, the T-dual type IIB magnetized D-brane scenario.

The presence of fluxes further complicates the constructions of these types of models: in the type IIB background, the $G_3$ fluxes give a large positive contribution to the D3-brane RR tadpoles, thus making it extremely hard to satisfy the D3 charge tadpole cancellation conditions by the magnetized D-brane sectors. In the first model with one unit of quantized flux [36], the large positive contribution to D3 charges from the flux, is cancelled by the "hidden sector" magnetized D9-branes, carrying negative D3 charges (first introduced for vacua without fluxes in [14]). Four-family models with one unit of flux and the starting SM-sector gauge symmetry (ii) were constructed by introducing a single stack of magnetized D9-branes with the negative D3 charge in [37]. These very few specific semirealistic constructions are extremely constrained, thus implying that semirealistic flux vacua are hard to come by. In this paper we advance this program in a new direction, by introducing *magnetized D9-branes with negative D3 charges as a part of the SM-sector* within frameworks (i) and (ii). As a consequence we obtain a large number of the three- and four-family SM-like flux with as much as three units of flux turned on. In particular *these constructions provide first four-dimensional $N = 1$ supersymmetric SM-like string vacua (i.e. three units of flux) as well as first examples of semirealistic SM-like string vacua with two units of flux* and many new models with one unit of flux. In addition, we also obtained SM-like flux models where the "hidden sector" D7-brane strong gauge dynamics in the IR region can definitely generate Veneziano-Yankielowicz superpotential, which is expected to be able to help stabilize the toroidal Kähler moduli by employing the KKLT mechanism [44].

In the following subsections we shall present explicit representative models within its class. Within each class there are typically more models and a sizable number of models within each class has been obtained by running computer code. The representative model in each class is typically chosen to have a minimal number of chiral (mainly right) exotics. In the following, we give a concise description of the representative models. (Please, refer to the appendix for tables containing these models and a detailed explanation of the notation employed.) A concise discussion of phenomenological implications of these models will be given in [47].

### A. Models with supersymmetric fluxes

In this subsection, we construct SM-like string vacua with the supersymmetric flux configuration, i.e. three units of quantized flux. Again, the key feature is the introduction of magnetized D9-branes with the negative D3 charge, which is a part of the SM-sector. These are the first three- and four-family SM-models within the supersymmetric flux background. The D-brane configurations of Model $- T_1 - 3$ and Model $- F_1 - 3$ are given in Table VII and Table XV, respectively. The chiral spectrum for the three-family model (Model $- T_1 - 3$) is given in Table II. All three toroidal Kähler moduli in these models are fixed by the supersymmetry conditions for the D-brane sector. [Note that the open string moduli and the Kähler moduli can form combined D-flat directions, corresponding to the brane recombination (see Ref. [2] in the context of specific type IIA orientifold compactifications). However, due to the flux backreaction it is expected that the open string moduli could become massive [30]; in this case the super-

TABLE II. The chiral spectrum in the open string sector of Model $- T_1 - 3$.

| Model $- T_1 - 3$ | $U(4)_C \times U(2)_L \times U(2)_R$ | $Q_4$ | $Q_{2L}$ | $Q_{2R}$ | $Q_{em}$ | $B-L$ | Field |
|---|---|---|---|---|---|---|---|
| $ab$ | $3 \times (4, \overline{2}, 1)$ | 1 | $-1$ | 0 | $-\frac{1}{3}, \frac{2}{3}, -1, 0$ | $\frac{1}{3}, -1$ | $Q_L, L_L$ |
| $ab'$ | $1 \times (\overline{4}, \overline{2}, 1)$ | $-1$ | $-1$ | 0 | $\frac{1}{3}, -\frac{2}{3}, 1, 0$ | $-\frac{1}{3}, 1$ | |
| $ac$ | $12 \times (\overline{4}, 1, 2)$ | $-1$ | 0 | 1 | $\frac{1}{3}, -\frac{2}{3}, 1, 0$ | $-\frac{1}{3}, 1$ | $Q_R, L_R$ |
| $ac'$ | $10 \times (4, 1, 2)$ | 1 | 0 | 1 | $-\frac{1}{3}, \frac{2}{3}, -1, 0$ | $\frac{1}{3}, -1$ | |
| $bc$ | $6 \times (1, \overline{2}, 2)$ | 0 | $-1$ | 1 | $-1, 0, 0, 1$ | 0 | $H$ |
| $bc'$ | $6 \times (1, 2, 2)$ | 0 | 1 | 1 | $-1, 0, 0, 1$ | 0 | $H$ |
| $b\square\!\square$ | $2 \times (1, 3, 1)$ | 0 | 2 | 0 | $0, \pm 1$ | 0 | |
| $b\begin{smallmatrix}\square\\\square\end{smallmatrix}$ | $2 \times (1, \overline{1}, 1)$ | 0 | $-2$ | 0 | 0 | 0 | |
| $c\square\!\square$ | $46 \times (1, 1, 3)$ | 0 | 0 | 2 | $0, \pm 1$ | 0 | |
| $c\begin{smallmatrix}\square\\\square\end{smallmatrix}$ | $146 \times (1, 1, 1)$ | 0 | 0 | 2 | 0 | 0 | |

symmetry conditions do stabilize Kähler moduli.] However, the SM Higgs doublets do not have Yukawa couplings to quarks and leptons due to the "wrong" quantum numbers under the global U(1) symmetries, and one has to look for new ways to generate quark and lepton masses. We shall further discuss the masses of the SM families as well as SM chiral exotics in [47].

### B. Models with nonsupersymmetric fluxes

In this subsection, we shall consider the string vacua with nonsupersymmetric fluxes, i.e. with two and one units of quantized fluxes. With fewer units of quantized flux, there is more freedom in satisfying the tadpole conditions. In the following we only present some typical three- and four-family models for each phenomenologically interesting case. Since the gauge symmetry breaking chain for each model can be determined easily from the analysis at the beginning of this section, we mainly focus on additional phenomenological aspects of these models.

#### *A. Two flux units models*

We constructed the first three- and four-family SM-like string vacua with two units of quantized flux with the representative models: Model $- T_1 - 2$ (Table VIII), Model $- F_1 - 2$ (Table XVI), Model $- F_2 - 2$ (Table XVII), and Model $- F_3 - 2$ (Table XVIII).

For the three-family model (Model $- T_1 - 2$), its chiral spectrum is given in Table III. Note that there is one left-chiral exotic $(\overline{4}, \overline{2}, 1, 1)$, which has Yukawa couplings to the right-chiral ones and SM Higgs doublet pairs (Higgs bi-doublets), thus these exotics can obtain a mass at the electroweak scale. In addition, there are five pairs of (nonchiral) SM Higgs doublet pairs, arising in the $bc$ sector when the $b$ and $c$ stacks of D-branes are coincident on the first two-torus; these SM Higgs doublet pairs have correct global U(1) quantum numbers, allowing for the Yukawa couplings to quarks and leptons. At the tree-level these Yukawa couplings can give masses and mixings for the second and third family fermions. The inclusion of quantum corrections may allow for the mass hierarchy and mixings for all three families. Note also that the three-family models with one unit of quantized flux, which have these similar nice features, can be constructed easily by choosing appropriate wrapping numbers for the $U(2)_R$ branes, e.g., $(-2, -1)(3, 1)(2, 1)$.

For the four-family models, there do not exist any left-chiral exotics. For the Model $- F_3 - 2$, we can give masses to one-family SM fermions. However, the SM fermion masses and mixings cannot be generated at the

TABLE III. The chiral spectrum in the open string sector of Model $- T_1 - 2$.

| Model $- T_1 - 2$ | $U(4)_C \times U(2)_L \times U(2)_R \times Sp(4)$ | $Q_4$ | $Q_{2L}$ | $Q_{2R}$ | $Q_{em}$ | $B-L$ | Field |
|---|---|---|---|---|---|---|---|
| $ab$ | $3 \times (4, \overline{2}, 1)$ | 1 | $-1$ | 0 | $-\frac{1}{3}, \frac{2}{3}, -1, 0$ | $\frac{1}{3}, -1$ | $Q_L, L_L$ |
| $ab'$ | $1 \times (\overline{4}, \overline{2}, 1)$ | $-1$ | $-1$ | 0 | $\frac{1}{3}, -\frac{2}{3}, 1, 0$ | $-\frac{1}{3}, 1$ | |
| $ac$ | $8 \times (\overline{4}, 1, 2)$ | $-1$ | 0 | 1 | $\frac{1}{3}, -\frac{2}{3}, 1, 0$ | $-\frac{1}{3}, 1$ | $Q_R, L_R$ |
| $ac'$ | $8 \times (4, 1, 2)$ | 1 | 0 | 1 | $-\frac{1}{3}, \frac{2}{3}, -1, 0$ | $\frac{1}{3}, -1$ | |
| $bc$ (Nonchiral) | $(1, 2, 2, 1) + (1, \overline{2}, \overline{2}, 1)$ | $\cdots$ | $\pm 1$ | $\pm 1$ | $\cdots$ | $\cdots$ | $H$ |
| $bc'$ | $4 \times (1, 2, 2)$ | 0 | 1 | 1 | $-1, 0, 0, 1$ | 0 | |
| $a(D7)_2$ | $1 \times (\overline{4}, 1, 1, 4)$ | $-1$ | 0 | 0 | $\frac{1}{6}, -\frac{1}{2}$ | $\frac{1}{3}, -1$ | |
| $b(D7)_2$ | $2 \times (1, 2, 1, 4)$ | 0 | 1 | 0 | $\pm\frac{1}{2}$ | 0 | |
| $c(D7)_2$ | $6 \times (1, 1, 2, 4)$ | 0 | 0 | 1 | $\pm\frac{1}{2}$ | 0 | |
| $c\square\!\square$ | $32 \times (1, 1, 3)$ | 0 | 0 | 2 | $0, \pm 1$ | 0 | |
| $c\begin{smallmatrix}\square\\\square\end{smallmatrix}$ | $112 \times (1, 1, 1)$ | 0 | 0 | 2 | 0 | 0 | |

tree level in the Model $- F_1 - 2$ and Model $- F_2 - 2$ because the SM Higgs doublet pairs have wrong quantum numbers under the U(1) global symmetries and thus no Yukawa couplings to quarks and leptons.

### *2. One flux unit models with* $\mathbf{U(2)}_{L,R}$ *negative D3 charge branes*

When one unit of flux is turned on, there is a wealth of models and these constructions of three- or four-family SM-like flux vacua can be classified in the following way:

(1) Model $- T_1 - 1$ (Table IX) and Model $- F_1 - 1$ (Table XIX). Except for some symmetric and antisymmetric representations, these two models do not contain any bi-fundamental chiral exotics in the observable sector. Their chiral spectra are given in Table IV and Table V, respectively. In particular, even though for the four-family model (Model $- F_1 - 1$) there is an anomaly-free $U(1)_R$ gauge symmetry, it is broken at the "right-handed" scale, when the $SU(2)_R \times U(1)_{B-L}$ gauge symmetry is broken down to the $U(1)_Y$ by giving VEVs to the scalar components of right-handed neutrino chiral superfields. However, in these two models, the SM Higgs doublet pairs do not have Yukawa couplings to quarks and leptons, due to the wrong global U(1) quantum numbers.

(2) In the Model $- T_2 - 1$ (Table X), Model $- T_3 - 1$ (Table XI) and Model $- F_3 - 1$ (Table XXI), there are four or five pairs of nonchiral Higgs bi-doublets arising from the $bc$ sector which can couple via Yukawa couplings to quarks and leptons, however, we can only give masses to one-family SM fermions. Also, some of the additional chiral exotics can obtain large masses by coupling to these nonchiral Higgs bi-doublets.

Even though these models employ nonchiral Higgs bi-doublets to give masses to fermions, we can easily find models with chiral Higgs bi-doublets with appropriate Yukawa couplings to quarks and leptons, see, e.g., Model $- T_6 - 1$ (Table XIV) and Model $- F_2 - 1$ (Table XX).

Another interesting four-family model is Model $- F_4 - 1$ (Table XXII); it does not contain any chiral exotics in the observable sector. The chiral Higgs bi-doublets in the $bc$ sector allow for the tree-level Yukawa couplings to quarks and leptons which can generate SM fermion masses and mixings for two families. Note also that the $U(4)_C$ symmetry emerges by Sp(16) D7-brane splitting at the string scale and thus $U(1)_C$ gauge symmetry is nonanomalous and is broken at the "right-handed" scale by the VEVs of the scalar components of the right-handed neutrino chiral superfields.

### *3. One flux unit models with* $\mathbf{U(4)}_C$ *negative D3 charge branes*

In this construction the $U(4)_C$ is due to the negative D3 charge magnetized D9-branes. Model $- T_4 - 1$ (Table XII) and Model $- F_5 - 1$ (Table XXIII) are such three- and four-family SM-like models. In both models, the $SU(2)_R$ gauge symmetry is generated by D7-brane

TABLE IV. The chiral spectrum in the open string sector of Model $- T_1 - 1$.

| Model $- T_1 - 1$ | $U(4)_C \times Sp(2)_L \times U(2)_R \times Sp(4)$ | $Q_4$ | $Q_{2R}$ | $Q_{em}$ | $B - L$ | Field |
|---|---|---|---|---|---|---|
| $ab$ | $3 \times (4, \overline{2}, 1, 1)$ | 1 | 0 | $-\frac{1}{3}, \frac{2}{3}, -1, 0$ | $\frac{1}{3}, -1$ | $Q_L, L_L$ |
| $ac$ | $3 \times (\overline{4}, 1, 2, 1)$ | $-1$ | 1 | $\frac{1}{3}, -\frac{2}{3}, 1, 0$ | $-\frac{1}{3}, 1$ | $Q_R, L_R$ |
| $bc$ | $8 \times (1, \overline{2}, 2, 1)$ | 0 | 1 | $-1, 0, 0, 1$ | 0 | $H$ |
| $c$(D3) | $1 \times (1, 1, \overline{2}, 4)$ | 0 | $-1$ | $\pm\frac{1}{2}$ | 0 | |
| $c\,\square\!\square$ | $23 \times (1, 1, 3, 1)$ | 0 | 2 | $0, \pm 1$ | 0 | |
| $c\,\begin{smallmatrix}\square\\\square\end{smallmatrix}$ | $73 \times (1, 1, 1, 1)$ | 0 | 2 | 0 | 0 | |

TABLE V. The chiral spectrum in the open string sector of Model $- F_1 - 1$.

| Model $- F_1 - 1$ | $U(4)_C \times U(2)_L \times Sp(8)_R \times Sp(8)$ | $Q_4$ | $Q_{2L}$ | $Q_{em}$ | $B - L$ | Field |
|---|---|---|---|---|---|---|
| $ab$ | $4 \times (4, \overline{2}, 1, 1)$ | 1 | $-1$ | $-\frac{1}{3}, \frac{2}{3}, -1, 0$ | $\frac{1}{3}, -1$ | $Q_L, L_L$ |
| $ac$ | $1 \times (\overline{4}, 1, 8, 1)$ | $-1$ | 0 | $\frac{1}{3}, -\frac{2}{3}, 1, 0$ | $-\frac{1}{3}, 1$ | $Q_R, L_R$ |
| $bc$ | $4 \times (1, \overline{2}, 8, 1)$ | 0 | $-1$ | $-1, 0, 0, 1$ | 0 | $H$ |
| $a(\mathrm{D7})_2$ | $2 \times (4, 1, 1, 8)$ | 1 | 0 | $\frac{1}{6}, -\frac{1}{2}$ | $\frac{1}{3}, -1$ | |
| $b(\mathrm{D7})_2$ | $4 \times (1, \overline{2}, 1, 8)$ | 0 | $-1$ | $\pm\frac{1}{2}$ | 0 | |
| $a\,\overline{\square\!\square}$ | $2 \times (\overline{10}, 1, 1, 1)$ | $-2$ | 0 | $\frac{1}{3}, -1$ | $\frac{2}{3}, -2$ | |
| $a\,\begin{smallmatrix}\square\\\square\end{smallmatrix}$ | $2 \times (6, 1, 1, 1)$ | 2 | 0 | $\frac{1}{3}, -1$ | $\frac{2}{3}, -2$ | |
| $b\,\overline{\square\!\square}$ | $10 \times (1, \overline{3}, 1, 1)$ | 0 | $-2$ | $0, \pm 1$ | 0 | |
| $b\,\begin{smallmatrix}\square\\\square\end{smallmatrix}$ | $54 \times (1, \overline{1}, 1, 1)$ | 0 | $-2$ | 0 | 0 | |

TABLE VI. The chiral spectrum in the open string sector of Model $- F_5 - 1$.

| Model $- F_5 - 1$ | $\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{Sp}(8)_R \times \mathrm{USp}(4)$ | $Q_4$ | $Q_{2L}$ | $Q_{em}$ | $B - L$ | Field |
|---|---|---|---|---|---|---|
| $ab$ | $4 \times (4, \overline{2}, 1, 1)$ | 1 | 0 | $-\frac{1}{3}, \frac{2}{3}, -1, 0$ | $\frac{1}{3}, -1$ | $Q_L, L_L$ |
| $ac$ | $1 \times (\overline{4}, 1, 8, 1)$ | $-1$ | 1 | $\frac{1}{3}, -\frac{2}{3}, 1, 0$ | $-\frac{1}{3}, 1$ | $Q_L, L_L$ |
| $bc$ (Nonchiral) | $(1, 2, 8, 1) + (1, \overline{2}, 8, 1)$ | $\cdots$ | $\pm 1$ | $\cdots$ | $\cdots$ | $H$ |
| $a(\mathrm{D7})_1$ | $4 \times (4, 1, 1, 4)$ | 1 | 0 | $\frac{1}{6}, -\frac{1}{2}$ | $\frac{1}{3}, -1$ | |
| $a_{\Box\Box}$ | $2 \times (10, 1, 1, 1)$ | 2 | 0 | $\frac{1}{3}, -1$ | $\frac{2}{3}, -2$ | |
| $a_{\overset{\Box}{\Box}}$ | $30 \times (6, 1, 1, 1)$ | 2 | 0 | $0, \pm 1$ | $\frac{2}{3}, -2$ | |
| $b_{\Box\Box}$ | $2 \times (1, 3, 1, 1)$ | 0 | 2 | $0, \pm 1$ | 0 | |
| $b_{\overset{\Box}{\Box}}$ | $2 \times (1, \overline{1}, 1, 1)$ | 0 | $-2$ | 0 | 0 | |

splitting, which yields eight and four copies of right-chiral representations, respectively. In particular, the four-family model (Model $- F_5 - 1$), whose chiral spectrum is given in Table VI, have several nice phenomenological features:

(i) No additional bi-fundamental chiral exotics in the observable sector;
(ii) Fermion masses and mixings for two families can be generated by the tree-level Yukawa couplings to the SM Higgs doublet pairs;
(iii) No additional electroweak scale U(1) symmetry.

### C. Models with Kähler moduli stabilized by D7-brane gauge dynamics

In this subsection, we consider the possibility of realizing a stabilization of the toroidal Kähler moduli stabilization à la KKLT mechanism [44]. In the original paper, the KKLT vacua are achieved via three steps:

(1) Turning on self-dual three-form fluxes on type IIB Calabi-Yau manifold; the flux-induced superpotential will fix the dilaton and all the complex structure moduli;
(2) Introducing a Kähler-moduli dependent nonperturbative superpotential, due to D7-brane strong infrared gauge dynamics or Euclidean D3-brane instanton effect. The vacuum is a supersymmetric anti de Sitter one with the Kähler-moduli fixed;
(3) Adding a set of anti-D3-branes to lift the anti de Sitter vacuum to a de Sitter one.

We only will focus on the second step, i.e. the generation of the nonperturbative superpotential, since it may help us stabilize all the toroidal moduli fields.

In our framework, the type IIB SM-like flux vacua typically require a stack of (filler) D3-branes, sitting on the $O3$-planes, in order to cancel the D3 charge tadpoles due to redundant negative D3 charges introduced by the magnetized D9-branes. This stack of D3-branes has typically a negative beta function; it thus possesses a nonperturbative gauge dynamics that results in a nonperturbative superpotential, due to gaugino condensation. However, this superpotential depends only on the dilaton-axion field and is independent Kähler-moduli. Thus, in order to generate the nonperturbative superpotential that depends on the toroidal Kähler moduli, the strong gauge dynamics has to arise due to D7-branes [44].

Recently, it was suggested [39] that the flux-induced soft mass terms may help decouple the open string moduli on D7-branes, leaving an infrared gauge theory with strong dynamics on the world-volume of this stack of D7-branes (in the hidden sector). However, in the concrete constructions of SM-like flux vacua a stack of hidden sector D7-branes typically do not have a negative beta function; this is due to the fact that the magnetized D9-branes with negative D3 charge have large "intersecting numbers" with the D7-branes and thus a large number of chiral matter, charged under D7-brane gauge symmetry. Additional chiral matter typically drastically modify the infrared gauge dynamics. [Light chiral matter can influence the (supersymmetric) gauge dynamics in two key aspects (see, e.g., [49] and references therein): (i) the chiral matter contributes to the beta function and thus affect the infrared dynamics (phase structure); (ii) it may induce matter condensation and contribute to the nonperturbative superpotential.]

TABLE VII. D-brane configurations and intersection numbers for Model $- T_1 - 3$.

| Model $- T_1 - 3$ | | $[\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{U}(2)_R]_{\mathrm{Observable}}$ | | | | | | | |
|---|---|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1, m^1)(n^2, m^2)(n^3, m^3)$ | $n_{\Box\Box}$ | $n_{\overset{\Box}{\Box}}$ | $b$ | $b'$ | $c$ | $c'$ | Kähler moduli |
| $a$ | 8 | $(1, 0)(1, 1)(1, -1)$ | 0 | 0 | $-3$ | 1 | 12 | $-10$ | $\chi_3 = \chi_2 = 2\chi_1$ |
| $b$ | 4 | $(1, 1)(2, -1)(1, 0)$ | $-2$ | 2 | $\cdots$ | $\cdots$ | 6 | $-6$ | $\chi_3 = 2\sqrt{10}$ |
| $c$ | 4 | $(-2, -1)(4, 1)(3, 1)$ | $-46$ | $-146$ | $\cdots$ | $\cdots$ | $\cdots$ | $\cdots$ | |

In this paper, we present the first SM-like flux vacua with strong gauge dynamics, resulting in gaugino condensation, on D7-branes of the hidden sector: Model $- T_5 - 1$ (Table XIII) and Model $- F_6 - 1$ (Table XXIV). These models have three- and four-family fermions (and additional chiral exotics), respectively. In these models the two (out of three) toroidal Kähler moduli have been fixed by supersymmetry conditions. For the four-family model (Model $- F_6 - 1$), its hidden sector is composed of two stacks of D7-branes denoted by $(D7)_1$ and $(D7)_2$, respectively. Both of them carry Sp(4) gauge symmetry and the associated beta functions are $-3(0)$ and $-5(-2)$, respectively. Here the beta functions in the brackets include the one-loop contribution from the open string moduli. These open string moduli are expected to become massive due to the flux backreaction (see, e.g.,[30]), and in this case the strong gauge dynamics can in principle induce the nonperturbative superpotential. Note, however, that gauge dynamics of the $(D7)_1$-branes results in the superconformal or Coulomb phase regime and the nonperturbative superpotential cannot be dynamically generated. As a result, only $(D7)_2$-branes can generate a Veneziano-Yankielowicz-type superpotential, induced by gaugino and matter condensations. A similar analysis can be applied to the three-family model (Model $- T_5 - 1$): in the case that open string moduli are decoupled [the beta function is $-9(0)$], the Veneziano-Yankielowicz superpotential is generated due to the gaugino and matter condensates on the D7-brane.

## IV. DISCUSSIONS AND CONCLUSIONS

In this paper we have advanced a program for explicit constructions of supersymmetric SM-like string vacua with supergravity fluxes turned on. In particular we obtained large classes of such type IIB models on $Z_2 \times Z_2$ orientifolds with one, two, and three units of quantized flux turned on. These models provide an important stepping stone toward broader classes of realistic constructions that not only contain the three- (or four-) family SM sector, but also stabilize some of the moduli (typically all toroidal moduli can be fixed).

Before our work, techniques for constructions of the chiral D-brane sectors flux vacua on type IIB orientifolds were developed in [30,31], however, the semirealistic constructions remained elusive until recently. The technical reason for such difficulties are large positive three-form flux contributions to the D3 charge in the internal space (D3-branes RR-tadpole), which makes the D3 charge conservation constraint hard to satisfy. In the first examples of SM flux vacua on $Z_2 \times Z_2$ orientifolds this constraint has been satisfied [36,37] by the introduction of the hidden sector magnetized D9-branes which carry negative D3 charges. (These types of D-branes were first introduced in the type IIA context with intersecting D6-branes without fluxes in [14].) Within this framework one three-family [36] and one four-family [37] SM-like models with one unit of quantized flux were obtained. However, one remains to be confronted by serious problems:

(i) There remains the outstanding problem of constructing semirealistic SM-like string vacua with supersymmetric fluxes, that for $Z_2 \times Z_2$ orientifolds corresponding to three units of quantized flux. The focus therefore shifted to the study of nonsupersymmetric flux effects as the key to break supersymmetry and a detailed study of soft supersymmetry breaking masses due to fluxes [39,40]. The flux-induced supersymmetry breaking, however, leads typically to soft masses $M_{\text{soft}} \sim \frac{M_s^2}{M_{\text{Pl}}}$, which implies an intermediate string scale or inhomogeneous warp factor in the internal space in order to stabilize the electroweak scale;

(ii) In spite of the successes in stabilizing the dilaton and toroidal complex structure moduli, the toroidal Kähler moduli are not completely fixed, thus one remains faced with the large vacuum degeneracy problem;

(iii) The fact that (supersymmetric) supergravity fluxes are abundant makes it imperative to extend the constructions to semirealistic models with more than one unit of quantized flux.

In this paper, we have made important progress in addressing the above issues. The key ingredient in the new SM-like flux constructions is the introduction of the negative D3 charge magnetized D9-brane as a part of the SM sector. These constructions turned out to be less constraining and resulted in three- and four-family SM-like string vacua with up to three units of quantized flux, thus leading to the first fully supersymmetric SM-like flux vacua. In addition to toroidal complex structure moduli and the dilaton being fixed by the flux, the supersymmetry conditions in the D-brane sector typically fix all the toroidal Kähler moduli, and the string scale is close to the Planck scale. We also constructed the first three- and four-family SM-like models with two units of the quantized flux and many new models with one unit of the quantized flux. Typically the representative models have (mainly right-) chiral exotics in the SM sector; however, we also have presented a few three- and four-family models with no SM-sector chiral exotics. (Note, the models do have additional tensor fields, i.e. chiral superfields in the symmetric and/or antisymmetric representation of the unitary gauge symmetry, typically associated with the negative D3 charge magnetized D9-brane.) We also have been able to construct SM-like flux vacua with strong infrared gauge dynamics on the hidden sector D7-branes, which leads to a nonperturbative superpotential that may fix the remaining toroidal Kähler modulus.

The SM Higgs doublet pairs appear at the intersections of the $SU(2)_L$ and $SU(2)_R$. For most models, SM fermion masses for one family, or SM fermion masses and mixings for two families can be generated by the tree-level Yukawa couplings. In the latter case the inclusion of quantum corrections may allow for the mass hierarchy and mixings of all three families. However, in few cases, most notably for the three-family SM-like model with the supersymmetric flux, all such Higgs fields have the wrong global U(1) quantum numbers that prevent them from coupling to quarks and leptons via Yukawa couplings; these models therefore face serious phenomenological difficulties. Note also that typically in these models some of the SM chiral exotics can obtain masses at least at the electroweak scale, due to the Yukawa couplings of the SM Higgs doublet pairs to such exotics.

In spite of a number of advances made with these new constructions, there are open problems which deserve further study. In particular, further discussion of the specific Yukawa couplings to quarks, leptons, and chiral exotics, as well as the subsequent further implications for the masses and mixings of the SM chiral fermions are needed [47]. For SM-like models with supersymmetric fluxes, one has to address the supersymmetry breaking mechanism. In principle the supersymmetry breaking could be due to the hidden D-brane strong gauge dynamics; however, the present models do not possess such a sector. Complete stabilization of all the moduli (and not only the toroidal closed sector one) remains an open problem. We postpone these issues for future research.

## ACKNOWLEDGMENTS

We would like to thank Paul Langacker for many discussions, encouragement, and collaboration on topics related to this paper. The research was supported in part by the National Science Foundation under Grant Nos. INT02-03585 (M. C.) and PHY-0070928 (T. Li), by the Department of Energy Grant DOE-EY-76-02-3071 (M. C., T. Liu), and by the Fay R. and Eugene L. Langberg Chair (M. C.).

## APPENDIX: D-BRANE CONFIGURATIONS AND INTERSECTION NUMBERS FOR SM-LIKE FLUX VACUA

Here we tabulate D-brane configurations and intersection numbers for the representative three- and four-family models within our new setup. In the first column of each table, $a$, $b$, and $c$ denote the U(4) (Sp(16)), $\mathrm{U}(2)_L$ [$\mathrm{Sp}(2)_L$ or $\mathrm{Sp}(4)_L$] and $\mathrm{U}(2)_R$ [$\mathrm{Sp}(4)_R$ or $\mathrm{Sp}(8)_R$] stacks of branes, respectively. D3, $(\mathrm{D7})_1$, $(\mathrm{D7})_2$, and $(\mathrm{D7})_3$ represent the filler branes along respective $\Omega R$, $\Omega R\omega$, $\Omega R\theta\omega$, and $\Omega R\theta$ orientifold planes, resulting in $\mathrm{Sp}(N)$ gauge groups. $N$, in the second column, corresponds to the number of filler D-branes in each stack. The third column depicts the wrapping numbers of the various D-branes. The intersection numbers between the various D-brane stacks are given in the remaining right columns where $b'$ and $c'$ are, respectively, the $\Omega R$ images of $b$ and $c$. In addition, the number of symmetric and antisymmetric chiral superfield representations for specific D-brane configurations is given. For convenience, we also tabulate the relations among the three toroidal Kähler moduli parameters $\chi_i$, imposed by the supersymmetry conditions. The model labels "$\mathrm{Model}-(T,F)_i-n$" appearing in the tables denote the "Model − (three, four family)$_i$− ($n$ units of fluxes)." [Since all the models have an even number of chiral supermultiplets in the fundamental representation of the $\mathrm{Sp}(N)$ gauge groups, these models are automatically free of discrete global gauge anomalies [50]. Finally, we emphasize that in this paper we do not fix the convention between the chirality and the sign of the intersection number. Instead, we consider the SU(2) D-branes that carry the more realistic chiral spectrum (typically only three- or four-families) as the $\mathrm{SU}(2)_L$ D-branes. These representative models therefore possess (mainly) right-chiral exotics.

TABLE VIII. D-brane configurations and intersection numbers for $\mathrm{Model}-T_1-2$.

| $\mathrm{Model}-T_1-2$ | | $[\mathrm{U}(4)_C\times\mathrm{U}(2)_L\times\mathrm{U}(2)_R]_{\mathrm{Observable}}\times[\mathrm{Sp}(4)]_{\mathrm{Hidden}}$ | | | | | | |
|---|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1,m^1)(n^2,m^2)(n^3,m^3)$ | $n_{\Box\Box}$ | $n_{\substack{\Box\\\Box}}$ | $b$ | $b'$ | $c$ | $c'$ |
| $a$ | 8 | $(1,0)(1,1)(1,-1)$ | 0 | 0 | $-3$ | 1 | 8 | $-8$ |
| $b$ | 4 | $(2,1)(2,-1)(1,0)$ | 0 | 0 | $\cdots$ | $\cdots$ | 0 | $-4$ |
| $c$ | 4 | $(-2,-1)(3,1)(3,1)$ | $-32$ | $-112$ | $\cdots$ | $\cdots$ | $\cdots$ | $\cdots$ |
| $(D7)_2$ | 4 | $(0,1)(1,0)(0,-1)$ | $\chi_3=\chi_2=\chi_1=\sqrt{21}$ | | | | | |

TABLE IX. D-brane configurations and intersection numbers for $\mathrm{Model}-T_1-1$.

| $\mathrm{Model}-T_1-1$ | | $[\mathrm{U}(4)_C\times\mathrm{Sp}(2)_L\times\mathrm{U}(2)_R]_{\mathrm{Observable}}\times[Sp(4)]_{\mathrm{Hidden}}$ | | | | | |
|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1,m^1)(n^2,m^2)(n^3,\tilde{m}^3)$ | $n_{\Box\Box}$ | $n_{\substack{\Box\\\Box}}$ | $b$ | $c$ | $c'$ |
| $a$ | 8 | $(1,0)(3,1)(3,-1/2)$ | 0 | 0 | $-3$ | 3 | 0 |
| $b$ | 2 | $(0,1)(0,-1)(2,0)$ | 0 | 0 | $\cdots$ | 8 | $\cdots$ |
| $c$ | 4 | $(-2,-1)(4,1)(3,1/2)$ | $-23$ | $-73$ | $\cdots$ | $\cdots$ | $\cdots$ |
| D3 | 4 | $(1,0)(1,0)(2,0)$ | $\chi_2=\chi_3,\ \frac{12}{\chi_2^2}+\frac{14}{\chi_1\chi_2}=1$ | | | | |

TABLE X. D-brane configurations and intersection numbers for Model $-\ T_2 - 1$.

| Model $-\ T_2 - 1$ | | $[\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{U}(2)_R]_{\mathrm{Observable}} \times [\mathrm{Sp}(4) \times \mathrm{Sp}(2)]_{\mathrm{Hidden}}$ | | | | | | |
|---|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1, m^1)(n^2, m^2)(n^3, \tilde{m}^3)$ | $n_{\Box\Box}$ | $n_{\substack{\Box\\\Box}}$ | $b$ | $b'$ | $c$ | $c'$ |
| $a$ | 8 | $(1, 0)(1, 1)(1, -1/2)$ | 0 | 0 | 3 | $-2$ | $-4$ | 4 |
| $b$ | 4 | $(2, -1)(1, 0)(5, 1/2)$ | 3 | $\cdots$ | $\cdots$ | $\cdots$ | 0 | $-4$ |
| $c$ | 4 | $(-2, 1)(3, -1)(3, -1/2)$ | 16 | 56 | $\cdots$ | $\cdots$ | $\cdots$ | $\cdots$ |
| D3 | 4 | $(1, 0)(1, 0)(2, 0)$ | | $\chi_3 = \chi_2 = \frac{5}{2}\chi_1 = \sqrt{39}$ | | | | |
| $(\mathrm{D7})_2$ | 2 | $(0, 1)(0, -1)(2, 0)$ | | | | | | |

TABLE XI. D-brane configurations and intersection numbers for Model $-\ T_3 - 1$.

| Model $-\ T_3 - 1$ | | $[\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{U}(2)_R]_{\mathrm{Observable}} \times [\mathrm{Sp}(8)]_{\mathrm{Hidden}}$ | | | | | | |
|---|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1, m^1)(n^2, m^2)(n^3, m^3)$ | $n_{\Box\Box}$ | $n_{\substack{\Box\\\Box}}$ | $b$ | $b'$ | $c$ | $c'$ |
| $a$ | 8 | $(1, 0)(1, 1)(1, -1)$ | 0 | 0 | 3 | 1 | $-3$ | 3 |
| $b$ | 4 | $(2, -1)(1, 0)(1, 2)$ | $-6$ | 6 | $\cdots$ | $\cdots$ | 0 | 12 |
| $c$ | 4 | $(-2, 1)(2, -1)(2, -1)$ | 10 | 54 | $\cdots$ | $\cdots$ | $\cdots$ | $\cdots$ |
| $(\mathrm{D7})_3$ | 8 | $(0, 1)(0, -1)(1, 0)$ | | $\chi_3 = \chi_2 = \frac{1}{4}\chi_1 = \sqrt{6}$ | | | | |

TABLE XII. D-brane configurations and intersection numbers for Model $-\ T_4 - 1$.

| Model $-\ T_4 - 1$ | | $[\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{Sp}(4)_R]_{\mathrm{Observable}} \times [Sp(4)]_{\mathrm{Hidden}}$ | | | | | |
|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1, m^1)(n^2, m^2)(n^3, m^3)$ | $n_{\Box\Box}$ | $n_{\substack{\Box\\\Box}}$ | $b$ | $b'$ | $c$ |
| $a$ | 8 | $(-1, -1)(2, 1)(2, 1)$ | $-2$ | $-30$ | 3 | $-5$ | $-4$ |
| $b$ | 4 | $(1, 0)(3, 1)(1, -1)$ | 4 | $-4$ | $\cdots$ | $\cdots$ | 0 |
| $c$ | 4 | $(1, 0)(0, 1)(0, -1)$ | 0 | 0 | $\cdots$ | $\cdots$ | $\cdots$ |
| D3 | 4 | $(1, 0)(1, 0)(1, 0)$ | | $3\chi_3 = \chi_2,\ \frac{12}{\chi_2^2} + \frac{8}{\chi_1\chi_2} = 1$ | | | |

TABLE XIII. D-brane configurations and intersection numbers for Model $-\ T_5 - 1$, here $\beta^g$ are beta functions for the associated gauge symmetries in the hidden sector.

| Model $-\ T_5 - 1$ | | $[\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{Sp}(8)_R]_{\mathrm{Observable}} \times [\mathrm{Sp}(8) \times \mathrm{Sp}(8)]_{\mathrm{Hidden}}$ | | | | | |
|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1, m^1)(n^2, m^2)(n^3, m^3)$ | $n_{\Box\Box}$ | $n_{\substack{\Box\\\Box}}$ | $b$ | $b'$ | $c$ |
| $a$ | 8 | $(1, 0)(1, 1)(1, -1)$ | 0 | 0 | 3 | $-3$ | $-1$ |
| $b$ | 4 | $(-2, -1)(2, 1)(2, 1)$ | $-10$ | $-54$ | $\cdots$ | $\cdots$ | $-4$ |
| $c$ | 8 | $(0, 1)(0, -1)(1, 0)$ | 0 | 0 | $\cdots$ | $\cdots$ | $\cdots$ |
| D3 | 8 | $(1, 0)(1, 0)(1, 0)$ | | $\chi_2 = \chi_3,\ \frac{4}{\chi_2^2} + \frac{8}{\chi_1\chi_2} = 1$ | | | |
| $(\mathrm{D7})_2$ | 8 | $(0, 1)(1, 0)(0, -1)$ | | $\beta^g_{\mathrm{D3}} = -14(-5),\ \beta^g_{(\mathrm{D7})_2} = -9(0)$ | | | |

TABLE XIV. D-brane configurations and intersection numbers for Model $-\ T_6 - 1$.

| Model $-\ T_6 - 1$ | | $[\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{U}(2)_R]_{\mathrm{Observable}} \times [\mathrm{Sp}(4) \times \mathrm{Sp}(4)]_{\mathrm{Hidden}}$ | | | | | | |
|---|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1, m^1)(n^2, m^2)(n^3, m^3)$ | $n_{\Box\Box}$ | $n_{\substack{\Box\\\Box}}$ | $b$ | $b'$ | $c$ | $c'$ |
| $a$ | 8 | $(1, 0)(1, 1)(1, -1)$ | 0 | 0 | 3 | $-1$ | $-6$ | 4 |
| $b$ | 4 | $(3, -1)(1, 0)(2, 1)$ | $-2$ | 2 | $\cdots$ | $\cdots$ | 5 | 5 |
| $c$ | 4 | $(-2, 1)(2, -1)(3, -1)$ | $-18$ | $-78$ | $\cdots$ | $\cdots$ | $\cdots$ | $\cdots$ |
| $(\mathrm{D7})_2$ | 8 | $(0, 1)(1, 0)(0, -1)$ | | $\chi_3 = \chi_2 = \frac{2}{3}\chi_1 = \sqrt{\frac{38}{3}}$ | | | | |
| $(\mathrm{D7})_3$ | 8 | $(0, 1)(0, -1)(1, 0)$ | | | | | | |

TABLE XV. D-brane configurations and intersection numbers for Model $-$ $F_1$ $-$ 3.

| Model $- F_1 - 3$ | $[\mathrm{U}(4)_C \times \mathrm{Sp}(4)_L \times \mathrm{U}(2)_R]_{\text{Observable}}$ | | | | | | | |
|---|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1, m^1)(n^2, m^2)(n^3, m^3)$ | $n_{\square\square}$ | $n_{\substack{\square\\\square}}$ | $b$ | $c$ | $c'$ | Kähler moduli |
| $a$ | 8 | $(1, 0)(2, 1)(1, -1)$ | 2 | $-2$ | $-2$ | 8 | $-12$ | $\chi_2 = 2\chi_3$ |
| $b$ | 4 | $(0, 1)(0, -1)(1, 0)$ | 0 | 0 | $\cdots$ | 8 | $\cdots$ | $\frac{24}{\chi_2^2} + \frac{20}{\chi_1\chi_2} = 1$ |
| $c$ | 4 | $(-2, -1)(4, 1)(3, 1)$ | $-46$ | $-146$ | $\cdots$ | $\cdots$ | $\cdots$ | |

TABLE XVI. D-brane configurations and intersection numbers for Model $-$ $F_1$ $-$ 2.

| Model $- F_1 - 2$ | $[\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{U}(2)_R]_{\text{Observable}} \times [\mathrm{Sp}(4)]_{\text{Hidden}}$ | | | | | | | |
|---|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1, m^1)(n^2, m^2)(n^3, m^3)$ | $n_{\square\square}$ | $n_{\substack{\square\\\square}}$ | $b$ | $b'$ | $c$ | $c'$ |
| $a$ | 8 | $(1, 0)(2, 1)(1, -1)$ | 2 | $-2$ | $-4$ | 0 | 4 | $-10$ |
| $b$ | 4 | $(1, 1)(2, -1)(1, 0)$ | $-2$ | 2 | $\cdots$ | $\cdots$ | 5 | $-3$ |
| $c$ | 4 | $(-2, -1)(3, 1)(3, 1)$ | $-32$ | $-112$ | $\cdots$ | $\cdots$ | $\cdots$ | $\cdots$ |
| $(\mathrm{D}7)_2$ | 4 | $(0, 1)(1, 0)(0, -1)$ | $\chi_2 = 2\chi_3 = 2\chi_1 = 3\sqrt{6}$ | | | | | |

TABLE XVII. D-brane configurations and intersection numbers for Model $-$ $F_2$ $-$ 2.

| Model $- F_2 - 2$ | $[\mathrm{U}(4)_C \times \mathrm{Sp}(4)_L \times \mathrm{U}(2)_R]_{\text{Observable}} \times [\mathrm{Sp}(8) \times \mathrm{Sp}(4)]_{\text{Hidden}}$ | | | | | | |
|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1, m^1)(n^2, m^2)(n^3, m^3)$ | $n_{\square\square}$ | $n_{\substack{\square\\\square}}$ | $b$ | $c$ | $c'$ |
| $a$ | 8 | $(1, 0)(2, 1)(1, -1)$ | 2 | $-2$ | $-2$ | 4 | $-10$ |
| $b$ | 4 | $(0, 1)(0, -1)(1, 0)$ | 0 | 0 | $\cdots$ | 6 | $\cdots$ |
| $c$ | 4 | $(-2, -1)(3, 1)(3, 1)$ | $-32$ | $-112$ | $\cdots$ | $\cdots$ | $\cdots$ |
| D3 | 8 | $(1, 0)(1, 0)(1, 0)$ | $2\chi_3 = \chi_2$ | | | | |
| $(\mathrm{D}7)_2$ | 4 | $(0, 1)(1, 0)(0, -1)$ | $\frac{18}{\chi_2^2} + \frac{18}{\chi_1\chi_2} = 1$ | | | | |

TABLE XVIII. D-brane configurations and intersection numbers for Model $-$ $F_3$ $-$ 2.

| Model $- F_3 - 2$ | $[\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{U}(2)_R]_{\text{Observable}} \times [Sp(4)]_{\text{Hidden}}$ | | | | | | | |
|---|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1, m^1)(n^2, m^2)(n^3, m^3)$ | $n_{\square\square}$ | $n_{\substack{\square\\\square}}$ | $b$ | $b'$ | $c$ | $c'$ |
| $a$ | 8 | $(1, 0)(2, 1)(1, -1)$ | 2 | $-2$ | $-3$ | $-1$ | 4 | $-10$ |
| $b$ | 4 | $(2, 1)(1, -1)(1, 0)$ | 2 | $-2$ | $\cdots$ | $\cdots$ | 0 | $-8$ |
| $c$ | 4 | $(-2, -1)(3, 1)(3, 1)$ | $-32$ | $-112$ | $\cdots$ | $\cdots$ | $\cdots$ | $\cdots$ |
| $(\mathrm{D}7)_2$ | 4 | $(0, 1)(1, 0)(0, -1)$ | $\chi_2 = 2\chi_3 = \frac{1}{2}\chi_1 = 3\sqrt{3}$ | | | | | |

TABLE XIX. D-brane configurations and intersection numbers for Model $-$ $F_1$ $-$ 1.

| Model $- F_1 - 1$ | $[\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{Sp}(8)_R]_{\text{Observable}} \times [Sp(8)]_{\text{Hidden}}$ | | | | | | |
|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1, m^1)(n^2, m^2)(n^3, m^3)$ | $n_{\square\square}$ | $n_{\substack{\square\\\square}}$ | $b$ | $b'$ | $c$ |
| $a$ | 8 | $(1, 0)(1, 1)(2, -1)$ | $-2$ | 2 | 4 | 0 | $-1$ |
| $b$ | 4 | $(-2, -1)(2, 1)(2, 1)$ | $-10$ | $-54$ | $\cdots$ | $\cdots$ | $-4$ |
| $c$ | 8 | $(0, 1)(0, -1)(1, 0)$ | 0 | 0 | $\cdots$ | $\cdots$ | $\cdots$ |
| $(\mathrm{D}7)_2$ | 8 | $(0, 1)(1, 0)(0, -1)$ | $\chi_3 = 2\chi_2, \frac{2}{\chi_2^2} + \frac{6}{\chi_1\chi_2} = 1$ | | | | |

TABLE XX. D-brane configurations and intersection numbers for Model $-\ F_2 - 1$.

| Model $-\ F_2 - 1$ | | $[\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{U}(2)_R]_{\text{Observable}} \times [\mathrm{Sp}(16) \times \mathrm{Sp}(4)]_{\text{Hidden}}$ | | | | | | |
|---|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1, m^1)(n^2, m^2)(n^3, m^3)$ | $n_{\Box\!\Box}$ | $n_{\substack{\Box\\ \Box}}$ | $b$ | $b'$ | $c$ | $c'$ |
| $a$ | 8 | $(1, 0)(1, 1)(2, -1)$ | $-2$ | 2 | 4 | 0 | $-4$ | 10 |
| $b$ | 4 | $(3, -1)(1, 0)(2, 1)$ | $-2$ | 2 | $\cdots$ | $\cdots$ | 5 | 5 |
| $c$ | 4 | $(-2, 1)(3, -1)(3, -1)$ | 32 | 112 | $\cdots$ | $\cdots$ | $\cdots$ | $\cdots$ |
| D3 | 16 | $(1, 0)(1, 0)(1, 0)$ | $\chi_3 = 2\chi_2 = \frac{2}{3}\chi_1 = \sqrt{30}$ | | | | | |
| $(\mathrm{D}7)_3$ | 4 | $(0, 1)(0, -1)(1, 0)$ | | | | | | |

TABLE XXI. D-brane configurations and intersection numbers for Model $-\ F_3 - 1$.

| Model $-\ F_3 - 1$ | | $[\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{U}(2)_R]_{\text{Observable}} \times [Sp(8)]_{\text{Hidden}}$ | | | | | | |
|---|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1, m^1)(n^2, m^2)(n^3, m^3)$ | $n_{\Box\!\Box}$ | $n_{\substack{\Box\\ \Box}}$ | $b$ | $b'$ | $c$ | $c'$ |
| $a$ | 8 | $(1, 0)(1, 1)(1, -1)$ | 0 | 0 | $-4$ | 2 | 4 | $-6$ |
| $b$ | 4 | $(2, 1)(3, -1)(1, 0)$ | $-2$ | 2 | $\cdots$ | $\cdots$ | 0 | 4 |
| $c$ | 4 | $(-2, -1)(2, 1)(3, 1)$ | $-18$ | $-78$ | $\cdots$ | $\cdots$ | $\cdots$ | $\cdots$ |
| $(\mathrm{D}7)_2$ | 8 | $(0, 1)(1, 0)(0, -1)$ | $\chi_3 = \chi_2 = \frac{3}{2}\chi_1 = \sqrt{21}$ | | | | | |

TABLE XXII. D-brane configurations and intersection numbers for Model $-\ F_4 - 1$.

| Model $-\ F_4 - 1$ | | $[\mathrm{Sp}(16)_C \times \mathrm{U}(2)_L \times \mathrm{U}(2)_R]_{\text{Observable}}$ | | | | | | | |
|---|---|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1, m^1)(n^2, m^2)(n^3, m^3)$ | $n_{\Box\!\Box}$ | $n_{\substack{\Box\\ \Box}}$ | $b$ | $b'$ | $c$ | $c'$ | Kähler moduli |
| $a$ | 16 | $(1, 0)(1, 0)(1, 0)$ | 0 | 0 | 1 | $\cdots$ | $-1$ | $\cdots$ | $\chi_1\chi_3 = 6$ |
| $b$ | 4 | $(2, -1)(0, 1)(3, -1)$ | $-10$ | 10 | $\cdots$ | $\cdots$ | 64 | 0 | $\frac{\chi_1}{\chi_2} + \frac{12}{\chi_1\chi_2} = 1$ |
| $c$ | 4 | $(-2, -1)(4, 1)(1, 1)$ | $-6$ | $-58$ | $\cdots$ | $\cdots$ | $\cdots$ | $\cdots$ | |

TABLE XXIII. D-brane configurations and intersection numbers for Model $-\ F_5 - 1$.

| Model $-\ F_5 - 1$ | | $[\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{Sp}(8)_R]_{\text{Observable}} \times [Sp(4)]_{\text{Hidden}}$ | | | | | |
|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1, m^1)(n^2, m^2)(n^3, m^3)$ | $n_{\Box\!\Box}$ | $n_{\substack{\Box\\ \Box}}$ | $b$ | $b'$ | $c$ |
| $a$ | 8 | $(-1, -1)(2, 1)(2, 1)$ | $-2$ | $-30$ | $-4$ | 0 | 1 |
| $b$ | 4 | $(1, 0)(1, 1)(2, -1)$ | $-2$ | 2 | $\cdots$ | $\cdots$ | 0 |
| $c$ | 8 | $(1, 0)(1, 0)(1, 0)$ | 0 | 0 | $\cdots$ | $\cdots$ | $\cdots$ |
| $(\mathrm{D}7)_1$ | 4 | $(1, 0)(0, 1)(0, -1)$ | $\chi_3 = 2\chi_2,\ \frac{2}{\chi_2^2} + \frac{3}{\chi_1\chi_2} = 1$ | | | | |

TABLE XXIV. D-brane configurations and intersection numbers for Model $-\ F_6 - 1$, here $\beta^g$ are beta functions for the associated gauge symmetries in the hidden sector.

| Model $-\ F_6 - 1$ | | $[\mathrm{U}(4)_C \times \mathrm{Sp}(8)_L \times \mathrm{U}(2)_R]_{\text{Observable}} \times [\mathrm{Sp}(4) \times \mathrm{Sp}(4)]_{\text{Hidden}}$ | | | | | |
|---|---|---|---|---|---|---|---|
| j | $N$ | $(n^1, m^1)(n^2, m^2)(n^3, m^3)$ | $n_{\Box\!\Box}$ | $n_{\substack{\Box\\ \Box}}$ | $b$ | $c$ | $c'$ |
| $a$ | 8 | $(1, 0)(1, 1)(1, -1)$ | 0 | 0 | $-1$ | 6 | $-4$ |
| $b$ | 8 | $(0, 1)(0, -1)(1, 0)$ | 0 | 0 | $\cdots$ | 3 | $\cdots$ |
| $c$ | 4 | $(-1, -1)(3, 1)(2, 1)$ | $-4$ | $-44$ | $\cdots$ | $\cdots$ | $\cdots$ |
| $(\mathrm{D}7)_1$ | 4 | $(1, 0)(0, 1)(0, -1)$ | $\chi_2 = \chi_3,\ \frac{6}{\chi_3^2} + \frac{5}{\chi_1\chi_3} = 1$ | | | | |
| $(\mathrm{D}7)_2$ | 4 | $(0, 1)(1, 0)(0, -1)$ | $\beta^g_{(\mathrm{D}7)_1} = -3(0),\ \beta^g_{(\mathrm{D}7)_2} = -5(-2)$ | | | | |

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