Physics Letters B 668 (2008) 63–66

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Realistic Type IIB supersymmetric Minkowski flux vacua

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ARTICLE INFO

Article history:
Received 6 May 2008
Received in revised form 3 August 2008
Accepted 6 August 2008
Available online 9 August 2008
Editor: M. Cvetič

PACS:
11.10.Kk
11.25.Mj
11.25.-w
12.60.Jv

ABSTRACT

We show that there exist supersymmetric Minkowski vacua on Type IIB toroidal orientifold with general flux compactifications where the RR tadpole cancellation conditions can be relaxed elegantly. Then we present a realistic Pati–Salam like model. At the string scale, the gauge symmetry can be broken down to the Standard Model (SM) gauge symmetry, the gauge coupling unification can be achieved naturally, and all the extra chiral exotic particles can be decoupled so that we have the supersymmetric SMs with/without SM singlet(s) below the string scale. The observed SM fermion masses and mixings can also be obtained. In addition, the unified gauge coupling, the dilaton, the complex structure moduli, the real parts of the Kähler moduli and the sum of the imaginary parts of the Kähler moduli can be determined as functions of the four-dimensional dilaton and fluxes, and can be estimated as well.

1. Introduction

One of the great challenging and essential problems in string phenomenology is the construction of the realistic string vacua, which can give us the low energy supersymmetric Standard Models (SMs) without exotic particles, and can stabilize the moduli fields. With renormalization group equation running, we can connect such constructions to the low energy realistic particle physics which will be tested at the upcoming Large Hadron Collider (LHC). During the last a few years, the intersecting D-brane models on Type II orientifolds [1], where the chiral fermions arise from the intersections of D-branes in the internal space [2] and the T-dual description in terms of magnetized D-branes [3], have been particularly interesting [4].

On Type IIA orientifolds with intersecting D6-branes, many non-supersymmetric three-family Standard-like models and Grand Unified Theories (GUTs) were constructed in the beginning [5]. However, there generically existed uncancelled Neveu–Schwarz–Neveu–Schwarz (NSNS) tadpoles and the gauge hierarchy problem. To solve these problems, semi-realistic supersymmetric Standard-like and GUT models have been constructed in Type IIA theory on the $\mathbf{T}^6/(\mathbb{Z}_2\times\mathbb{Z}_2)$ orientifold [6,7] and other backgrounds [8]. Interestingly, only the Pati–Salam like models can give all the Yukawa couplings. Without the flux background, Pati–Salam like models have been constructed systematically in Type IIA theory on the $\mathbf{T}^6/(\mathbb{Z}_2\times\mathbb{Z}_2)$ orientifold [7]. Although we may explain the SM fermion masses and mixings in one model [9], the moduli fields have not been stabilized, and it is very difficult to decouple the chiral exotic particles. To stabilize the moduli via supergravity fluxes, the flux models on Type II orientifolds have also been constructed [10–14]. Especially, for the supersymmetric AdS vacua on Type IIA orientifolds with flux compactifications, the RR tadpole cancellation conditions can be relaxed [13,14]. And then we can construct the flux models that can explain the SM fermion masses and mixings [14]. However, those models are in the AdS vacua and have quite a few chiral exotic particles that are difficult to be decoupled.

In this Letter, we consider the Type IIB toroidal orientifold with the Ramond–Ramond (RR), NSNS, non-geometric and S-dual flux compactifications [15]. We find that the RR tadpole cancellation conditions can be relaxed elegantly in the supersymmetric Minkowski vacua, and then we may construct the realistic Pati–Salam like models [16]. In this Letter, we present a concrete simple model which is very interesting from the phenomenological point of view and might describe Nature. We emphasize that we do not fix the four-dimensional dilaton via flux potential, and our model is a solution to the equations of motion for all the Type IIB fields.

2. Type IIB flux compactifications

We consider the Type IIB string theory compactified on a $\mathbf{T}^6$ orientifold where $\mathbf{T}^6$ is a six-torus factorized as $\mathbf{T}^6=\mathbf{T}^2\times\mathbf{T}^2\times\mathbf{T}^2$ whose complex coordinates are z_i, $i=1,2,3$ for the ith two-torus,

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doi:10.1016/j.physletb.2008.08.010

Table 1
General spectrum for magnetized D-branes on the Type IIB $\mathbf{T}^6$ orientifold

Sector	Representation
aa	$U(N_a)$ vector multiplet
	3 adjoint multiplets
$ab+ba$	I_{ab} $(N_a, \bar{N}_b)$ multiplets
$ab'+b'a$	$I_{ab'}$ (N_a, N_b) multiplets
$aa'+a'a$	$\frac{1}{2}(I_{aa'} - I_{aO3})$ symmetric multiplets
	$\frac{1}{2}(I_{aa'} + I_{aO3})$ anti-symmetric multiplets

respectively. The orientifold projection is implemented by gauging the symmetry ΩR, where Ω is world-sheet parity, and R is given by

$$R : (z_1, z_2, z_3) \to (-z_1, -z_2, -z_3). \tag{1}$$

Thus, the model contains 64 O3-planes. In order to cancel the negative RR charges from these O3-planes, we introduce the magnetized D$(3+2n)$-branes which are filling up the four-dimensional Minkowski space–time and wrapping $2n$-cycles on the compact manifold. Concretely, for one stack of N_a D-branes wrapped m_a^i times on the ith two-torus $\mathbf{T}_i^2$, we turn on n_a^i units of magnetic fluxes F_a^i for the center of mass $U(1)_a$ gauge factor on $\mathbf{T}_i^2$, such that

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

where m_a^i can be half integer for tilted two-torus. Then, the D9-, D7-, D5- and D3-branes contain 0, 1, 2 and 3 vanishing m_a^is, respectively. Introducing for the ith two-torus the even homology classes $[\mathbf{0}_i]$ and $[\mathbf{T}_i^2]$ for the point and two-torus, respectively, the vectors of the RR charges of the ath stack of D-branes and its image are

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

respectively. The "intersection numbers" in Type IIA language, which determine the chiral massless spectrum, are

$$I_{ab} = [\Pi_a] \cdot [\Pi_b] = \prod_{i=1}^{3} \big(n_a^i m_b^i - n_b^i m_a^i\big). \tag{4}$$

Moreover, for a stack of N D$(2n+3)$-branes whose homology classes on $\mathbf{T}^6$ is (not) invariant under ΩR, we obtain a $USp(2N)$ $(U(N))$ gauge symmetry with three anti-symmetric (adjoint) chiral superfields due to the orbifold projection. The physical spectrum is presented in Table 1.

The flux models on Type IIB orientifolds with four-dimensional $N=1$ supersymmetry are primarily constrained by the RR tadpole cancellation conditions that will be given later, the four-dimensional $N=1$ supersymmetric D-brane configurations, and the K-theory anomaly free conditions. For the D-branes with world-volume magnetic field $F_a^i = n_a^i/(m_a^i \chi_i)$ where χ_i is the area of $\mathbf{T}_i^2$ in string units, the condition for the four-dimensional $N=1$ supersymmetric D-brane configurations is

$$\sum_i \big(\tan^{-1}\big(F_a^i\big)^{-1} + \theta\big(n_a^i\big)\pi\big) = 0 \bmod 2\pi, \tag{5}$$

where $\theta(n_a^i) = 1$ for $n_a^i < 0$ and $\theta(n_a^i) = 0$ for $n_a^i \geqslant 0$. The K-theory anomaly free conditions are [17]

$$\sum_a N_a m_a^1 m_a^2 m_a^3 = \sum_a N_a m_a^1 n_a^2 n_a^3 = \sum_a N_a n_a^1 m_a^2 n_a^3$$
$$= \sum_a N_a n_a^1 n_a^2 m_a^3 = 0 \bmod 2. \tag{6}$$

And the holomorphic gauge kinetic function for a generic stack of D$(2n+3)$-branes is given by [16,18,19]

$$f_a = \frac{1}{\kappa_a}\big(n_a^1 n_a^2 n_a^3 s - n_a^1 m_a^2 m_a^3 t_1 - n_a^2 m_a^1 m_a^3 t_2 - n_a^3 m_a^1 m_a^2 t_3\big), \tag{7}$$

where κ_a is equal to 1 and 2 for $U(n)$ and $USp(2n)$, respectively.

We turn on the NSNS flux h_0, RR flux e_i, non-geometric fluxes b_{ii} and $\bar{b}_{ii}$, and the S-dual fluxes f_i, g_{ij} and g_{ii} [15]. To avoid the subtleties, these fluxes should be even integers due to the Dirac quantization. For simplicity, we assume

$$e_i = e, \qquad b_{ii} = \beta, \qquad \bar{b}_{ii} = \bar{\beta},$$
$$f_i = f, \qquad g_{ij} = -g_{ii} = g, \tag{8}$$

where $i \neq j$. Then the constraint on fluxes from Bianchi identities is

$$f\bar{\beta} = g\beta. \tag{9}$$

The RR tadpole cancellation conditions are

$$\sum_a N_a n_a^1 n_a^2 n_a^3 = 16, \qquad \sum_a N_a n_a^i m_a^j m_a^k = -\frac{1}{2} e\bar{\beta},$$
$$N_{\mathrm{NS7}_i} = 0, \qquad N_{\mathrm{I7}_i} = 0, \tag{10}$$

where $i \neq j \neq k \neq i$, and the N_{NS7_i} and N_{I7_i} denote the NS7 brane charge and the other 7-brane charge, respectively [15]. Thus, if $e\bar{\beta} < 0$, the RR tadpole cancellation conditions are relaxed elegantly because $-e\bar{\beta}/2$ only needs to be even integer. Moreover, we have 7 moduli fields in the supergravity theory basis, the dilaton s, three Kähler moduli t_i, and three complex structure moduli u_i. With the above fluxes, we can assume

$$t \equiv t_1 + t_2 + t_3, \qquad u_1 = u_2 = u_3 \equiv u. \tag{11}$$

Then the superpotential becomes

$$\mathcal{W} = 3ieu + ih_0 s - t\big(\beta u - i\bar{\beta}u^2\big) - st(f - igu). \tag{12}$$

The Kähler potential for these moduli is

$$\mathcal{K} = -\ln(s + \bar{s}) - \sum_{i=1}^{3} \ln(t_i + \bar{t}_i) - \sum_{i=1}^{3} \ln(u_i + \bar{u}_i). \tag{13}$$

In addition, the supergravity scalar potential is

$$V = e^{\mathcal{K}}\big(\mathcal{K}^{i\bar{j}} D_i\mathcal{W} D_{\bar{j}}\mathcal{W} - 3|\mathcal{W}|^2\big), \tag{14}$$

where $\mathcal{K}^{i\bar{j}}$ is the inverse metric of $\mathcal{K}_{i\bar{j}} \equiv \partial_i\partial_{\bar{j}}\mathcal{K}$, $D_i\mathcal{W} = \partial_i\mathcal{W} + (\partial_i\mathcal{K})\mathcal{W}$, and $\partial_i = \partial_{\phi_i}$ where ϕ_i can be s, t_i, and u_i. Thus, for the supersymmetric Minkowski vacua, we have

$$\mathcal{W} = \partial_s\mathcal{W} = \partial_t\mathcal{W} = \partial_u\mathcal{W} = 0. \tag{15}$$

From $\partial_s\mathcal{W} = \partial_t\mathcal{W} = 0$, we obtain

$$t = \frac{ih_0}{f - igu}, \qquad s = -\frac{\beta}{f}u, \tag{16}$$

then the superpotential turns out

$$\mathcal{W} = \left(3e - \frac{h_0\beta}{f}\right)iu. \tag{17}$$

Therefore, to satisfy $\mathcal{W} = \partial_u\mathcal{W} = 0$, we obtain

$$3ef = \beta h_0. \tag{18}$$