Physics Letters B 665 (2008) 267–270

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

A realistic world from intersecting D6-branes

Ching-Ming Chen [a], Tianjun Li [a,b,*], V.E. Mayes [a], Dimitri V. Nanopoulos [a,c,d]

[a] George P. and Cynthia W. Mitchell Institute for Fundamental Physics, Texas A&M University, College Station, TX 77843, USA
[b] Institute of Theoretical Physics, Chinese Academy of Sciences, Beijing 100080, China
[c] Astroparticle Physics Group, Houston Advanced Research Center (HARC), Mitchell Campus, Woodlands, TX 77381, USA
[d] Academy of Athens, Division of Natural Sciences, 28 Panepistimiou Avenue, Athens 10679, Greece

ARTICLE INFO

Article history:
Received 17 January 2008
Received in revised form 11 June 2008
Accepted 13 June 2008
Available online 19 June 2008
Editor: M. Cvetič

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

ABSTRACT

We describe a three-family Pati–Salam model from intersecting D6-branes in type IIA string theory on the $\mathbf{T}^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold which is of strong phenomenological interest. In the model, the gauge coupling unification is achieved naturally at the string scale, and the gauge symmetry can be broken down to the Standard Model (SM) close to the string scale. Moreover, we find that it is possible to obtain the correct SM quark masses and mixings, and the tau lepton mass. Additionally, neutrino masses and mixings may be generated via the seesaw mechanism. Furthermore, we calculate the supersymmetry breaking soft terms, and the corresponding low-energy supersymmetric particle spectra which may potentially be tested at the Large Hadron Collider (LHC), and provide the observed dark matter density.

Published by Elsevier B.V.

1. Introduction

Although string theory has long teased us with her power to encompass all known physical phenomena in a complete mathematical structure, an actual worked out example of observed physics is still lacking. Indeed, the major problem of string phenomenology is to construct *at least one* realistic model with all moduli stabilized, which completely describes known physics as well as potentially being predictive of unknown phenomena. With the dawn of the LHC era, new discoveries will hopefully be upon us. In particular, supersymmetry is expected to be found as well as the Higgs states required to break the electroweak symmetry. Therefore, it would be highly desirable to have a complete model derived from string theory which is able to make predictions for the supersymmetric particle and Higgs spectra, as well as describing currently known particle physics. In this Letter, we are embarking on such an enterprise.

During the last few years, 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 shown great promise in model building [4–7]. The appeal of intersecting D-brane models has been in part based upon the fact that chiral fermions are present at the intersections of different stacks of branes and the multiplicity of such fermions is given by the topologically invariant intersection number. However, there are two serious problems in almost all supersymmetric D-brane models: the absence of gauge coupling unification at the string scale, and the rank one problem in the Standard Model (SM) fermion Yukawa matrices. Thus, a comprehensive phenomenological study of a concrete model from the string scale to the weak scale has yet to be made (for the previous phenomenology study, please see Refs. [8,9]). Therefore, the first major problem is whether or not it is possible to have a supersymmetric intersecting D-brane model which might describe Nature at some point(s) or subspace(s) of its moduli space. Following this, we can then consider the other problems, for example, the moduli stabilization and fine-tuning problems, etc. Interestingly, for the first time we find that there is a single intersecting D6-brane model on the type IIA $\mathbf{T}^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold where the above problems can be solved [6,10]. In this model, we will show that the gauge coupling unification can be realized naturally at string scale, the realistic Yukawa mass matrices can be generated, and the realistic superpartner spectra with a relic neutralino density of phenomenological interest can be obtained. We emphasize that we will not consider the moduli stabilization in this Letter since it is not our goal here. The moduli stabilization problem is very important, and will be studied in the similar model building on type IIB orientifolds with general flux compactifications [11,12].

2. Model building

The model [6,10] is constructed in type IIA string theory compactified on a $\mathbf{T}^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold where $\mathbf{T}^6$ is a six-torus

* Corresponding author at: George P. and Cynthia W. Mitchell Institute for Fundamental Physics, Texas A&M University, College Station, TX 77843, USA.
E-mail address: junlt@physics.tamu.edu (T. Li).

0370-2693/$ – see front matter Published by Elsevier B.V.
doi:10.1016/j.physletb.2008.06.024

Table 1
D6-brane configurations and intersection numbers where (n^i, l^i) are the wrapping numbers for the one cycle on the ith two torus

		$U(4)_C \times U(2)_L \times U(2)_R \times USp(2)^4$										
	N	$(n^1, l^1) \times (n^2, l^2) \times (n^3, l^3)$	n_S	n_A	b	b'	c	c'	1	2	3	4
a	8	$(0,-1) \times (1,1) \times (1,1)$	0	0	3	0	−3	0	1	−1	0	0
b	4	$(3,1) \times (1,0) \times (1,-1)$	2	−2	–	–	0	0	0	1	0	−3
c	4	$(3,-1) \times (0,1) \times (1,-1)$	−2	2	–	–	–	–	−1	0	3	0
1	2	$(1,0) \times (1,0) \times (2,0)$	$\chi_1 = 3,\ \chi_2 = 1,\ \chi_3 = 2$									
2	2	$(1,0) \times (0,-1) \times (0,2)$	$\beta_1^g = -3,\ \beta_2^g = -3$									
3	2	$(0,-1) \times (1,0) \times (0,2)$	$\beta_3^g = -3,\ \beta_4^g = -3$									
4	2	$(0,-1) \times (0,1) \times (2,0)$										

Table 2
The chiral and vector-like superfields, and their quantum numbers under the gauge symmetry $SU(4)_C \times SU(2)_L \times SU(2)_R \times USp(2)_1 \times USp(2)_2 \times USp(2)_3 \times USp(2)_4$

	Quantum number	Q_4	Q_{2L}	Q_{2R}	Field
ab	$3 \times (4, \bar{2}, 1, 1, 1, 1, 1)$	1	−1	0	$F_L(Q_L, L_L)$
ac	$3 \times (\bar{4}, 1, 2, 1, 1, 1, 1)$	−1	0	1	$F_R(Q_R, L_R)$
$a1$	$1 \times (4, 1, 1, 2, 1, 1, 1)$	1	0	0	
$a2$	$1 \times (\bar{4}, 1, 1, 1, 2, 1, 1)$	−1	0	0	
$b2$	$1 \times (1, 2, 1, 1, 2, 1, 1)$	0	1	0	
$b4$	$3 \times (1, \bar{2}, 1, 1, 1, 1, 2)$	0	−1	0	
$c1$	$1 \times (1, 1, \bar{2}, 2, 1, 1, 1)$	0	0	−1	
$c3$	$3 \times (1, 1, 2, 1, 1, 2, 1)$	0	0	1	
b_S	$2 \times (1, 3, 1, 1, 1, 1, 1)$	0	2	0	T_L^i
b_A	$2 \times (1, \bar{1}, 1, 1, 1, 1, 1)$	0	−2	0	S_L^i
c_S	$2 \times (1, 1, \bar{3}, 1, 1, 1, 1)$	0	0	−2	T_R^i
c_A	$2 \times (1, 1, 1, 1, 1, 1, 1)$	0	0	2	S_R^i
ab'	$3 \times (4, 2, 1, 1, 1, 1, 1)$	1	1	0	
	$3 \times (\bar{4}, \bar{2}, 1, 1, 1, 1, 1)$	−1	−1	0	
ac'	$3 \times (4, 1, 2, 1, 1, 1, 1)$	1		1	Φ_i
	$3 \times (\bar{4}, 1, \bar{2}, 1, 1, 1, 1)$	−1	0	−1	$\bar{\Phi}_i$
bc	$6 \times (1, 2, \bar{2}, 1, 1, 1, 1)$	0	1	−1	H_u^i, H_d^i
	$6 \times (1, \bar{2}, 2, 1, 1, 1, 1)$	0	−1	1	

factorized as $\mathbf{T}^6 = \mathbf{T}^2 \times \mathbf{T}^2 \times \mathbf{T}^2$ and the D6-branes wrap an one cycle on each two torus [5]. We present its D6-brane configurations and intersection numbers of the model in Table 1, and the resulting spectrum in Table 2 [6,10]. We put the a', b, and c stacks of D6-branes on the top of each other on the third two torus, and as a result there are additional vector-like particles from $N = 2$ subsectors.

The anomalies from three global $U(1)$s of $U(4)_C$, $U(2)_L$ and $U(2)_R$ are cancelled by the Green–Schwarz mechanism, and the gauge fields of these $U(1)$s obtain masses via the linear $B \wedge F$ couplings. Thus, the effective gauge symmetry is $SU(4)_C \times SU(2)_L \times SU(2)_R$. In order to break the gauge symmetry, on the first torus, we split the a stack of D6-branes into a_1 and a_2 stacks with 6 and 2 D6-branes, respectively, and split the c stack of D6-branes into c_1 and c_2 stacks with two D6-branes for each one. In this way, the gauge symmetry is further broken to $SU(3)_C \times SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$. Moreover, the $U(1)_{I_{3R}} \times U(1)_{B-L}$ gauge symmetry may be broken to $U(1)_Y$ by giving vacuum expectation values (VEVs) to the vector-like particles with the quantum numbers $(\mathbf{1}, \mathbf{1}, \mathbf{1/2}, -\mathbf{1})$ and $(\mathbf{1}, \mathbf{1}, -\mathbf{1/2}, \mathbf{1})$ under the $SU(3)_C \times SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$ gauge symmetry from $a_2 c_1'$ intersections [6,10].

Since the gauge couplings in the Minimal Supersymmetric Standard Model (MSSM) are unified at the GUT scale $\sim 2.4 \times 10^{16}$ GeV, the additional exotic particles present in the model must necessarily become superheavy. To accomplish this it is first assumed that the $USp(2)_1$ and $USp(2)_2$ stacks of D6-branes lie on the top of each other on the first torus, so we have two pairs of vector-like particles with $USp(2)_1 \times USp(2)_2$ quantum numbers $(2, 2)$. These particles can break $USp(2)_1 \times USp(2)_2$ down to the diagonal $USp(2)_{D12}$ near the string scale, and then states arising from intersections $a1$ and $a2$ may obtain vector-like masses close to the string scale. Moreover, we assume that the T_R^i and S_R^i obtain VEVs near the string scale, and their VEVs satisfy the D-flatness of $U(1)_R$. To preserve the D-flatness of $U(1)_L$, we assume that the VEVs of S_L^i is TeV scale. With T_R^i and S_R^i, we can give the GUT-scale masses to the particles from the intersections $c1$ and $c3$ via three point functions, and to T_R^i via four point functions [13]. However, the chiral exotic particles from the intersections $b2$ and $b4$ can be decoupled only at the intermediate scale, around 10^{12} GeV due to the strong dynamics in the hidden sector [13]. Also, one can show that the open string moduli (adjoint chiral supermultiplets) are exact flat directions to all orders in perturbation and thus cannot get masses and decouple. The non-perturbative mechanism and/or background fluxes may be able to give masses to these open string moduli. Moreover, to have one pair of light Higgs doublets, it is necessary to fine-tune the mixing parameters of the Higgs doublets. In particular, the μ term and the right-handed neutrino masses may be generated via the following high-dimensional operators that we introduce as in the effective field theory

$$W \supset \frac{y_\mu^{ijkl}}{M_{\rm St}} S_L^i S_R^j H_u^k H_d^l + \frac{y_{Nij}^{mnkl}}{M_{\rm St}^3} T_R^m T_R^n \Phi_i \Phi_j F_R^k F_R^l, \tag{1}$$

where y_μ^{ijkl} and y_{Nij}^{mnkl} are Yukawa couplings, and $M_{\rm St}$ is the string scale. Thus, the μ term is TeV scale and the right-handed neutrino masses can be in the range 10^{10-14} GeV for $y_\mu^{ijkl} \sim 1$ and $y_{Nij}^{mnkl} \sim 10^{(-7)-(-3)}$. Note that for the similar Pati–Salam model on type IIB orientifold with general flux compactifications, we can easily decouple all the chiral exotic particles [11].

3. Phenomenological consequences

In the string theory basis, we have the dilaton S, three Kähler moduli T^i, and three complex structure moduli U^i [14]. The U^i for the present model are

$$U^1 = 3i, \qquad U^2 = i, \qquad U^3 = -1 + i. \tag{2}$$

The corresponding moduli s, t^i and u^i in the supergravity theory basis are related to the S, T^i and U^i moduli by [14]

$$\begin{aligned} &{\rm Re}(s) = \frac{e^{-\phi_4}}{2\pi}\left(\frac{\sqrt{U_2^1 U_2^2 U_2^3}}{|U^1 U^2 U^3|}\right), \qquad {\rm Re}\big(t^j\big) = \frac{i\alpha'}{T^j}, \\ &{\rm Re}\big(u^j\big) = \frac{e^{-\phi_4}}{2\pi}\left(\sqrt{\frac{U_2^j}{U_2^k U_2^l}}\right)\left|\frac{U^k U^l}{U^j}\right|, \end{aligned} \tag{3}$$

where ϕ_4 is the four-dimensional dilaton, U_2^i is the imaginary part of U^i, and $j \neq k \neq l \neq j$.

The holomorphic gauge kinetic function for a generic P stack of D6-branes which does not lie on one of O6-planes, is given by [14]

$$f_P = \frac{1}{8}\big(2n_P^1 n_P^2 n_P^3 s - n_P^1 l_P^2 l_P^3 u^1 - n_P^2 l_P^1 l_P^3 u^2 - 2n_P^3 l_P^1 l_P^2 u^3\big). \tag{4}$$

And then we have

$$g_{SU(4)_C}^2 = g_{SU(2)_L}^2 = g_{SU(2)_R}^2 = \left[\frac{\sqrt{6} e^{-\phi_4}}{8\pi}\right]. \tag{5}$$

Thus, the gauge couplings for $SU(4)_C$, $SU(2)_L$ and $SU(2)_R$ in our model are unified at the string scale naturally. As in the Georgi–Glashow $SU(5)$ model, the Pati–Salam model has canonical $U(1)_Y$ normalization as well. So we have the canonical gauge coupling unification in our model. For simplicity, we neglect the little hierarchy between the string scale and the GUT scale, which may

Table 3
Low energy supersymmetric particles and their masses (in GeV)

h^0	H^0	A^0	$H^\pm$	$\tilde{g}$	$\chi_1^\pm$	$\chi_2^\pm$	χ_1^0	χ_2^0
117.5	907.4	907.4	911.3	2192	216.5	1229	197.1	216.5
χ_3^0	χ_4^0	$\tilde{t}_1$	$\tilde{t}_2$	$\tilde{u}_1/\tilde{c}_1$	$\tilde{u}_2/\tilde{c}_2$	$\tilde{b}_1$	$\tilde{b}_2$	
−1227	1228	1636	1965	2142	1945	1811	1968	
$\tilde{d}_1/\tilde{s}_1$	$\tilde{d}_2/\tilde{s}_2$	$\tilde{\tau}_1$	$\tilde{\tau}_2$	$\tilde{\nu}_\tau$	$\tilde{e}_1/\tilde{\mu}_1$	$\tilde{e}_2/\tilde{\mu}_2$	$\tilde{\nu}_e/\tilde{\nu}_\mu$	
2144	1944	253.9	1010	1002	1060	549.9	1056	

be explained via threshold corrections. Assuming the value of the unified gauge coupling in the MSSM, we obtain

$$e^{-\phi_4} = 20.1. \qquad (6)$$

Thus, the string scale is $\sim 2.1\times 10^{17}$ GeV for $M_{\rm St} = \pi^{1/2} e^{\phi_4} M_{\rm Pl}$ where $M_{\rm Pl}$ is the reduced Planck scale.

The Kähler metric for the chiral superfields from the intersections of the P and Q stacks of D6-branes is [14]

$$\tilde{K} \supset e^{\phi_4+\gamma_E\sum_{i=1}^3 \theta_{PQ}^i} \prod_{j=1}^{3}\left[\sqrt{\frac{\Gamma(1-\theta_{PQ}^i)}{\Gamma(\theta_{PQ}^i)}}\left(t^j+\bar{t}^j\right)^{-\theta_{PQ}^i}\right],$$

where γ_E is the Euler–Mascheroni constant, and θ_{PQ}^i is the suitable positive angle between the P and Q stacks of D6-branes on the ith two torus in units of π [13], and can be written as a function of s, u^i, and the wrapping numbers for the P and Q stacks of D6-branes.

The Kähler metric for the vector-like chiral superfields from the intersections of the P and Q stacks of D6-branes that are parallel on the jth two torus and intersect on the kth and lth two tori is given by [14]

$$\tilde{K} \supset \left[(s+\bar{s})\left(u^j+\bar{u}^j\right)\left(t^k+\bar{t}^k\right)\left(t^l+\bar{t}^l\right)\right]^{-1/2}. \qquad (7)$$

We emphasize that the problem of finding a mechanism leading to the low scale supersymmetry breaking in a natural and controlled way is an interesting and important question, but it is out of the scope of the present work. For simplicity, we assume that only the F terms of the complex structure moduli u^i break supersymmetry at the TeV scale and are parametrized as follows

$$F^{u^i} = \sqrt{3} m_{3/2}\left(u^i+\bar{u}^i\right)\Theta_i, \quad \text{for } i=1,2,3, \qquad (8)$$

where $m_{3/2}$ is the gravitino mass, and Θ_i are real numbers and satisfy $\sum_{i=1}^3 |\Theta_i|^2 = 1$. Then, we can calculate the gaugino masses (M_i), the universal scalar masses m_{FL} and m_{FR} respectively for the left-handed and right-handed SM fermions, the universal scalar mass m_H for Higgs fields H_u^i and H_d^i, and the universal trilinear soft term A_Y at the string scale [15]. Using the code `SuSpect` [16] and `MicrOMEGAs` [17], we can calculate the low energy supersymmetric particle spectrum and the dark matter density, respectively. With $m_{3/2} = 1100$ GeV, ${\rm Re}\, t_1 = 1/6.6$, ${\rm Re}\, t_2 = {\rm Re}\, t_3 = 0.5$, $\tan\beta = 46$, $m_{\rm top} = 170.9$ GeV, and positive μ and Θ_3, we show the neutralino dark matter relic density in the Θ_1–Θ_2 plane in Fig. 1 where the region with Higgs boson mass larger than 114 GeV is given as well. Therefore, we do have the parameter space that satisfies all the known experimental constraints and can give large enough dark matter density. As an example, for $\Theta_1 = -0.6$ and $\Theta_2 = 0.293$, we present the low energy supersymmetric particle spectrum in Table 3 which can be tested at the LHC, and we obtain the corresponding dark matter density $\Omega h^2 = 0.105$ which is very close to the observed value.

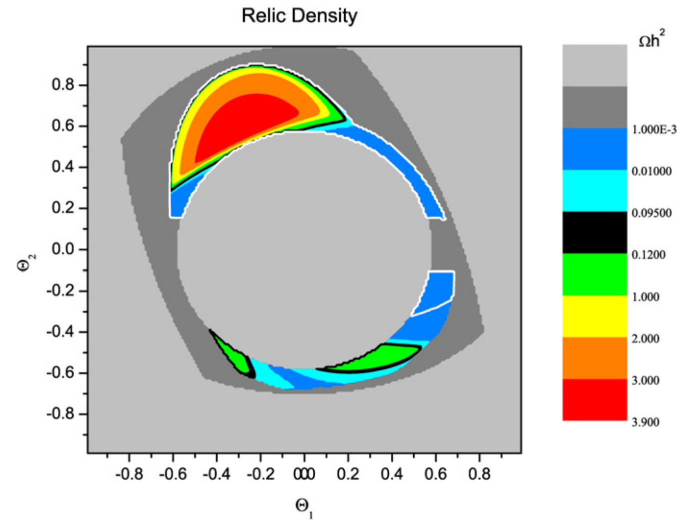

Fig. 1. Contour map of the neutralino dark matter relic density as a function of the goldstino angles Θ_1 and Θ_2 for $\tan\beta = 46$ with $\Theta_3 > 0$. The dark bands correspond to regions of the parameter space with the observed dark matter density while areas within the white contour denote regions where the Higgs mass is above the LEP limit, $m_h \geqslant 114$ GeV. The dark gray regions indicate regions where the neutralino is not LSP or other mass limits are not satisfied. The light gray regions are excluded by constraints on the soft terms at high scale.

4. The SM fermion masses and mixings

Because all the SM fermions and Higgs fields arise from the intersections on the first torus, we will only consider it for simplicity. The up-type quark mass matrix M^U at the GUT scale is [18]

$$c_0^U \begin{pmatrix} A^U H_u^1 + E^U H_u^4 & B^U H_u^3 + F^U H_u^6 & D^U H_u^2 + C^U H_u^5 \\ C^U H_u^3 + D^U H_u^6 & A^U H_u^5 + E^U H_u^2 & B^U H_u^1 + F^U H_u^4 \\ F^U H_u^2 + B^U H_u^5 & C^U H_u^1 + D^U H_u^4 & A^U H_u^3 + E^U H_u^6 \end{pmatrix}$$

where c_0^U is a constant which includes the quantum corrections and the contributions to the Yukawa couplings from the second and third two tori. The theta functions A^U, B^U, C^U, D^U, E^U, and F^U are

$$A^U \equiv \vartheta\begin{bmatrix}\epsilon^{U1}\\ \phi^{(1)}\end{bmatrix}\left(\kappa^{(1)}\right), \qquad B^U \equiv \vartheta\begin{bmatrix}\epsilon^{U1}+\frac{1}{3}\\ \phi^{(1)}\end{bmatrix}\left(\kappa^{(1)}\right),$$
$$C^U \equiv \vartheta\begin{bmatrix}\epsilon^{U1}-\frac{1}{3}\\ \phi^{(1)}\end{bmatrix}\left(\kappa^{(1)}\right), \qquad D^U \equiv \vartheta\begin{bmatrix}\epsilon^{U1}+\frac{1}{6}\\ \phi^{(1)}\end{bmatrix}\left(\kappa^{(1)}\right),$$
$$E^U \equiv \vartheta\begin{bmatrix}\epsilon^{U1}+\frac{1}{2}\\ \phi^{(1)}\end{bmatrix}\left(\kappa^{(1)}\right), \qquad F^U \equiv \vartheta\begin{bmatrix}\epsilon^{U1}-\frac{1}{6}\\ \phi^{(1)}\end{bmatrix}\left(\kappa^{(1)}\right),$$

where

$$\epsilon^{U1} \equiv \frac{\epsilon_c^{U1}-\epsilon_b^{U1}-2\epsilon_a^{U1}}{6}, \qquad \kappa^{(1)} \equiv \frac{6J^{(1)}}{\alpha'},$$
$$\phi^{(1)} = \theta_c^{(1)} - \theta_b^{(1)} - 2\theta_a^{(1)}, \qquad (9)$$

where ϵ_a^{U1}, ϵ_b^{U1} and ϵ_c^{U1} respectively are the shifts of a, b, and c stacks of D6-branes, $J^{(1)}$ is the Kähler modulus, and $\theta_a^{(1)}$, $\theta_b^{(1)}$ and $\theta_c^{(1)}$ are the Wilson line phases for the a, b, and c stacks on the first two torus, respectively.

At the GUT scale, the down-type quark mass matrix M^D is obtained from the above up-type quark mass matrix M^U by changing the upper index U and lower index u to D and d, respectively. The lepton mass matrix M^L is obtained from M^D by changing the upper index D to L. In addition, we emphasize that c_0^U, c_0^D, and c_0^L are constants and not matrices since the intersections ab and ac are 1 on the second and third two tori. Also, on the second two torus, we can split the a stack of D6-branes into a_1' and a_2' stacks with 6 and 2 D6-branes, respectively, and split the c stack of D6-branes into c'_1 and c'_2 stacks with two D6-branes for each one. And then, we obtain that c_0^U, c_0^D, and c_0^L can be different real numbers.

To generate the suitable SM fermion masses and mixings at the GUT scale, we choose $\epsilon^{U1} = \epsilon^{L1} = 0$, $\epsilon^{D1} = 0.061$, and $\kappa^{(1)} = 39.6i$. One pair of the Higgs doublets at low energy is fine-tuned to be

$$\begin{aligned} H_u &= 0.000187283 H_u^1 + 0.166161 H_u^2 + 0.703369 H_u^3 \\ &\quad + 0.690696 H_u^4 + 0.00338659 H_u^5 + 0.0242905 H_u^6, \\ H_d &= 0.00141716 H_d^1 + 0.999603 H_d^3 + 0.0281534 H_d^5 \\ &\quad + 6.3266 \times 10^{-5} H_d^6. \end{aligned} \tag{10}$$

Then, with suitable c_0^U, c_0^D, and c_0^L by adjusting the areas (triangles) on the second two torus, we obtain the SM fermion mass matrices at the GUT scale

$$M^U \simeq m_t \begin{pmatrix} 0.000266 & 0.00109 & 0.00747 \\ 0.00109 & 0.00481 & 0.0310 \\ 0.00747 & 0.0310 & 0.999 \end{pmatrix},$$

$$M^D \simeq m_b \begin{pmatrix} 0.00141 & 0.000025 & 4 \times 10^{-6} \\ 0.000155 & 0.028 & 0.0 \\ 0.0 & 2.2 \times 10^{-7} & 1 \end{pmatrix},$$

$$M^L \simeq m_\tau \begin{pmatrix} 0.00142 & 3.0 \times 10^{-6} & 2.8 \times 10^{-8} \\ 3.0 \times 10^{-6} & 0.0282 & 1.4 \times 10^{-9} \\ 2.8 \times 10^{-8} & 1.4 \times 10^{-9} & 1 \end{pmatrix}.$$

The above mass matrices can produce the correct quark masses and CKM mixings, and the correct τ lepton mass at the electroweak scale [19]. The electron mass is about 6.5 times larger that the expected value, while the muon mass is about 40% smaller. Similar to the GUTs [20], we have roughly the wrong fermion mass relation $m_e/m_\mu \simeq m_d/m_s$. In principle, the correct electron and muon masses can be generated via high-dimensional operators via the four point functions [21]. Moreover, neutrino masses and mixings can be generated via the seesaw mechanism by choosing suitable Majorana mass matrix for the right-handed neutrinos.

In short, in order to obtain the realistic supersymmetric particle spectra and explain the SM fermion masses and mixings, we have considered 15 independent free parameters: two supersymmetry F-term breaking terms Θ_1 and Θ_2, one shift ϵ^{D1}, one Kähler modulus $J^{(1)}$ (or $\kappa^{(1)}$), five relative VEVs for H_u^i, three relative VEVs for H_d^i, and three overall constants c_0^U, c_0^D, and c_0^L. Here, we do not count the extra overall scale factors for the SM fermion Yukawa couplings, and the parameters that are chosen to be zero in our numerical calculations.

5. Conclusions

We have briefly described a three-family intersecting D6-brane model where gauge coupling unification is achieved at the string scale and where the gauge symmetry can be broken to the Standard Model. In the model, it is possible to calculate the supersymmetry breaking soft terms and obtain the low energy supersymmetric particle spectrum within the reach of the LHC. Finally, it is possible to obtain the SM quark masses and CKM mixings and the tau lepton mass, and the neutrino masses and mixings may be generated via the seesaw mechanism. Although we have chosen specific values for the moduli fields to obtain agreement with experiments, it may be possible to uniquely predict these values by introducing the most general fluxes, which is under investigation.

Acknowledgements

This research was supported in part by the Mitchell-Heep Chair in High Energy Physics (C.M.C), by the Cambridge-Mitchell Collaboration in Theoretical Cosmology (T.L.), and by the DOE grant DE-FG03-95-Er-40917 (D.V.N.).

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