$E^2$ Instanton Effects and Higgs Physics In Intersecting Brane Models

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Abstract

String instanton effects in Higgs physics are discussed through a type IIA model based on $T^6/(Z^2 \times Z^2)$ orientifold compactification. By inclusion of rigid $E^2$-branes, the model exhibits an MSSM-like spectrum, as well as extra $\mu$ and quartic Higgs couplings. These extra couplings are induced via $E^2$ instantons non-perturbatively. Setting the string scale at $10^{18}$ GeV, one gets interesting TeV Higgs physics. In particular, the tree-level Higgs mass can be uplifted substantially.
1 Introduction

Recently, string instanton effects have been intensively explored in moduli stabilization of flux induced compactifications [1, 2, 3] (and references therein) and string phenomenology, especially for non-perturbative generation of right-handed neutrino masses and $\mu$ term in intersecting brane models [4, 5, 6, 7]. These non-perturbative effects come from nonzero global charges $Q_a = N_a \Xi \circ (\Pi_a - \Pi'_a)$ carried by the instantons, which lead to interesting charged matter couplings [8] (for recent reviews, see for example, [9, 10]). Setting the string scale at the order of $10^{18}$ GeV, one finds $m_\nu$ and $\mu$ in acceptable ranges without any fine-tuning.

In addition to the $\mu$ term, there is another important coupling in Higgs physics, the quartic coupling, which controls Higgs boson masses. In the minimal supersymmetric standard model (MSSM), the tree-level mass of the lightest Higgs particle $h$ is well below the LEPII bound. To make ends meet, one needs substantial radiative contribution to $m_h$ which is dominated by the stop quark [15, 19]. In order to obtain the desired uplifting, both the stop mass and the mixing have to be large. And this greatly constrains the parameter space in the MSSM and aggravates fine tuning problems associated with soft mass terms. This provides motivations to make extensions beyond the MSSM, such as the next leading-order minimal supersymmetric standard model (NMSSM) [14] and beyond minimal supersymmetric standard model (BMSSM) [12, 13]. In certain examples, extra quartic Higgs couplings are present which modify tree-level Higgs masses. Their significance is controlled by the mechanism of supersymmetry breaking in hidden sector and the value of the associated mass scale.

Motivated by the rich phenomenologies generated by stringy instantons, in this paper we will discuss their effects on two important mass scales in Higgs physics, i.e, the $\mu$ term and the mass scale $M$ associated with the quartic couplings, in $T^6/(Z^2 \times Z'^2)$ orientifold compactification of type IIA theories [20]. They are induced non-perturbatively via $E2$ instantons. Setting the string GUT scale at $10^{18}$ GeV, one gets interesting TeV Higgs physics. In particular, the tree-level Higgs mass can be uplifted substantially.

In section 2, a $\mathcal{N} = 1$ supersymmetric model is constructed that exhibits an MSSM-like spectrum (including the right-handed neutrino) with suitable wrapping numbers of $D6$ and $E2$ branes. In section 3, we discuss the generations of $E2$-branes induced $\mu$ term and quartic couplings. The structure of these quartic terms are explicitly calculated. They obviously modify the Higgs masses, which are expressed as expansions of a small
parameter \( \varepsilon \sim \mu/M \), as shown in section 4. We conclude in section 5.

## 2 The setup

We discuss an intersecting \( D6 \)-branes model in \( T^6/(Z^2 \times Z') \) orientifold of type IIA theories. All the moduli are stabilized if non-perturbative \( E2 \)-brane instanton effects are taken into account \[1\,2\,3\], and standard model spectrum can be obtained by properly arranging the intersecting branes. Shown in table 1 are the wrapping numbers of four-stack branes \( a, b, c, d \). The model carries gauge groups \( U(3)_a \times U(2)_b \times U(1)_c \times U(1)_d \), of which all the \( U(1)_i \) become massive by the Green-Schwarz mechanism except \( U(1)_Y \),

\[
Q_Y = \frac{1}{6}Q_a - \frac{1}{2}Q_c - \frac{1}{2}Q_d
\]

(2.1)

The gauge groups then conforms to that of MSSM-like theories. The intersecting number \( I_{cd} = -3 \) implies neutrinos \( \nu_R \) are also encoded. Shown in table 2 are the chiral spectra of theories corresponding to wrapping numbers in table 1.

For the model to be supersymmetric, each stack of branes has to satisfy two conditions \[16\],

\[
m_x^1 m_x^2 m_x^3 - \sum_{I \neq J \neq K} n_x^I n_x^J n_x^K U^I U^J = 0
\]

(2.2)

and

\[
n_x^1 n_x^2 n_x^3 - \sum_{I \neq J \neq K} m_x^I m_x^J n_x^K U^I U^J > 0
\]

(2.3)

where \( U^I = R^I_Y / R^I_X \) is the complex structure modulus of \( I \)th torus with radii \( R^I_X, R^I_Y \).

Note that in table 1, \( N_h D6 \)-branes and \( N_O O6 \) branes are added to cancel the tadpoles,

\[
\sum_{a=1}^{K} N_a (\Pi_a + \Pi'_a) = N_O \Pi_{O6}
\]

(2.4)

Also, stacks \( a \) and \( d \) are parallel in the transverse directions. The open string modes stretching between them are massive, of the order \( L/(\sqrt{2\pi\alpha_s}) \) (\( L \) is the transverse distance). So matter contents in table 2 are exact in the effective theory below the string scale. In addition, two \( E2 \)-branes \( M, N \) are embedded. We will see in the next section that they non-perturbatively induce interesting small \( \mu \) term and quartic terms in Higgs physics, respectively.
\begin{tabular}{|c|c|c|c|}
\hline
\(N_i\) & \((n_i^1, m_i^1)\) & \((n_i^2, m_i^2)\) & \((n_i^3, m_i^3)\) \\
\hline
\(N_a = 6\) & (1, 0) & (3, 1) & (3, -1/2) \\
\(N_b = 4\) & (1, 1) & (1, 0) & (1, -1/2) \\
\(N_c = 2\) & (0, 1) & (0, -1) & (2, 0) \\
\(N_d = 2\) & (1, 0) & (3, 1) & (3, -1/2) \\
\(N_h = 4\) & (-2, 1) & (-3, 1) & (-3, 1/2) \\
\(N_O = 6\) & (1, 0) & (1, 0) & (1, 0) \\
\(E2M\) & (1, 0) & (1, -1) & (1, 1/2) \\
\(E2N\) & \((n_N^1, -n_N^1)\) & \((n_N^2, \frac{12n_N^3}{1-s})\) & \((\frac{6n_N^3(n_N^3)^3}{1-s}, \frac{1}{n_N^3})\) \\
\hline
\end{tabular}

Table 1: Wrapping numbers of D6-branes and E2-instantons which wrap on a rigid three-cycle on \(Z^2 \times Z^2\) toroidal orientifold. \(n_N^1, n_N^2\) are real numbers \((s = (n_N^1)^2(n_N^2)^4)\). The model is supersymmetric if \(U_3 = 2U_1 = -2U_2 = 1\).

\begin{table}[h]
\begin{tabular}{|c|c|c|}
\hline
intersection numbers & matter & Rep \\
\hline
\(I_{ab} = I_{ab}^* = 3\) & \(Q_L\) & 3(3, 2) \\
\(I_{ac} = -3\) & \(U_R\) & 3(3, 1) \\
\(I_{ac^*} = 3\) & \(D_R\) & 3(3, 1) \\
\(I_{db} = I_{db^*} = 3\) & \(L\) & 3(1, 2) \\
\(I_{cd} = -3\) & \(\nu_R\) & 3(1, 1) \\
\(I_{cd^*} = 3\) & \(E_R\) & 3(1, 1) \\
\(I_{bc} = -1\) & \(H_u\) & 1(1, 2) \\
\(I_{be^*} = -1\) & \(H_v\) & 1(1, 2) \\
\hline
\end{tabular}
\end{table}

Table 2: Chiral matters spectrum for the wrapping numbers in table 1.
3 Non-perturbative Higgs physics from $E_2$ instanton

To yield a non-perturbative $\mu$ term, one assigns the following intersection numbers between $E_2$-brane and $D6_{b,c}$-branes

$$I_{Mb} = -1, \quad I_{Mb^*} = 0, \quad I_{Mc} = I_{Mc^*} = 1 \quad (I_{bc} < 0) \quad (3.1)$$

The intersection number $I_{Ma}$ also has to satisfy,

$$I_{Ma} - I_{Ma^*} = 0, \quad (\alpha = a, d) \quad (3.2)$$

in order to exclude the extra charged zero modes. The wrapping numbers on $E2_M$ are $2(1,0)(1,-1)(1,1/2)$, which are determined by the constraints Eq. (3.1) and Eq. (3.2), as shown in table 1. The number of triangles on each torus is 1, contributing to $H^i_u \lambda^i_a \bar{\lambda}^i_b e^{-A_i}$ and $H^i_d \lambda^i_a \bar{\lambda}^i_b e^{-A_i}$ terms respectively for intersecting $(b,c)$ and $(b,c^*)$ branes. This generates a $\mu H_u H_d$ term non-perturbatively in four-dimensional effective theory, as desired $[4, 6, 7]$.

We now discuss the quartic operator $\lambda^i_H (H_u H_d)^2$ and its implication for Higgs physics. These operators were constructed in certain BMSSM examples. They can greatly uplift Higgs masses when $M$ is in the range of $1 \sim 10$ TeV. Similar to the stringy instanton induced $\mu$ term as shown above, it is possible to construct these quartic terms non-perturbatively. That is, the roles played by hidden sectors to generate these operators in other models can be totally replaced by stringy instanton effect in our model.

In order to exclude extra zero modes on $D6_{a,d}$-branes, one has the constraints on the intersection number $E2_N$ and $D6$-branes

$$I_{Na} = I_{Na^*}, \quad (\alpha = a, d) \quad (3.3)$$

and

$$I_{Nb} = -4, \quad I_{Nb^*} = 0, \quad I_{Nc} = I_{Nc^*} = 2 \quad (I_{bc} < 0) \quad (3.4)$$

which can be obtained by counting the numbers of charged zero modes that arise from strings streching between the $E2_N$ and $D6_{b,c}$-branes.

As shown in table 2, the wrapping numbers of $E2_N$-brane are represented by two integer $(n^1_N, n^2_N)$. $E2_N$ also preserve the same supersymmetry as $D6$-branes, i.e, the wrapping numbers of $E2_N$ satisfy the constraints Eq. (2.2) and Eq. (2.3).
Figure 1: The left and right diagrams correspond to triangles on the first and second tori, respectively. In the first torus, $E_2$ and $c$ intersect twice, $A_1, A_2$ represent their areas. On the second and third tori, they intersect only once, whose areas are represented by $A_3$ and $A_4$.

The general strategy to compute charged matters coupling in $E_2$ instanton background has been outlined in [8]. In our case,

$$
< (H_u H_d)^2 >_{E_2-\text{inst}} = \int d^4x \sum_{\text{conf}} \prod_i^4 (d\lambda_{ai}) \times \prod_j^4 (d\bar{\lambda}_{aj}) e^{-S_{\text{inst}}} e^{Z'}
$$

$$
\times < H_u >_{\lambda_{a1}, \bar{\lambda}_{a1}} < H_d >_{\lambda_{a2}, \bar{\lambda}_{a2}} < H_u >_{\lambda_{a3}, \bar{\lambda}_{a3}} < H_d >_{\lambda_{a4}, \bar{\lambda}_{a4}} (3.5)
$$

which can be computed via conformal field theory techniques (see also [17, 18]). To appreciate the structure of Eq. (3.5), we take for example $n^1_N = 2, n^2_N = 1$. They are shown by three simple triangles in figure 1. Non-perturbative terms in each torus are proportional to

$$
(H_u^{ij} \lambda^i_a \bar{\lambda}^j_b e^{-A_1} + H_d^{ij} \lambda^i_a \bar{\lambda}^j_b e^{-A_1}) + (H_u^{ij} \lambda^i_a \bar{\lambda}^j_b e^{-A_2} + H_d^{ij} \lambda^i_a \bar{\lambda}^j_b e^{-A_2}),
$$

$$
H_u^{ij} \lambda^i_a \bar{\lambda}^j_b e^{-A_3} + H_d^{ij} \lambda^i_a \bar{\lambda}^j_b e^{-A_3},
$$

$$
H_u^{ij} \lambda^i_a \bar{\lambda}^j_b e^{-A_4} + H_d^{ij} \lambda^i_a \bar{\lambda}^j_b e^{-A_4},
$$

respectively. $A_i$ is the area in string units of the triangle as shown in the figure 1. Note that $A_2 = 4A_1$. The mixing terms between $H_u H_d$s are highly suppressed due to simplicities

\footnotetext[1]{Other choices of $n^1_N$ and $n^2_N$ will yield more complex expressions, but similar physics.}
of triangle structure on the second and third tori. This leads to the following term in the four-dimensional effective action,

\[ S_{\text{nonpert}} = \frac{A}{4!M} \varepsilon_{ijkl} \varepsilon_{mnpq} H_u^{lm} H_u^{jn} H_v^{kp} H_v^{lq}, \] (3.7)

where

\[ A = \frac{\pi^3}{4} (\Gamma_1 + \theta_{E2b,1} - \theta_{E2c,1} - \theta_{E2}) \sum_{i,j=1}^3 e^{-2(\tilde{A}_i + \tilde{A}_j)} \] (3.8)

and

\[ M = g_s M_s V_{E2} e^{S_{\text{inst}}(E2N)} \] (3.9)

where \( V_{E2} = \text{Vol}(E2)/l_3^3 \), \( \tilde{A}_i = A_i \), \( (i \neq 1, 2) \) and \( \tilde{A}_{1,2} = In(e^{A_1} + e^{A_2}) \). The rescaling for charged zero modes \( \lambda \rightarrow \lambda \sqrt{\frac{\pi}{g_s}} \) and the \( g_s \) factor independence for each disc imply that each disc diagram carries an overall normalization factor \( 2\pi/g_s \) \([5, 17]\). Thus, one gets

\[ \mu \sim g_s^{-1} M_s e^{-S_{\text{inst}}(E2M)} \] (3.10)

Eqs. (3.9) and (3.10) determine the significance of non-perturbative stringy effects on Higgs physics. With \( M_s \sim 10^{18} \) GeV and \( V_{E2M} \sim V_{E2N} \sim 10^{-30} \), one has \( \mu \sim 100 \) GeV and \( M \sim 1 \) TeV. The \( V_{E2M,N} \) values will increase as \( n_{1N}^1 \) and \( n_{1N}^1 \) decrease (the ratio of \( \text{Vol}_{E2N}/\text{Vol}_{E2M} \) is smaller). Without any fine tuning, these mass scales are exactly in the range desired by phenomenology. We will see in the next section, in particular, tree-level Higgs masses can be greatly uplifted.

There are other possible \( E2N \) instanton induced charged matter couplings. Note that for \( I_{Nh} \neq 0 \) and \( I_{hc} = I_{hc'} = 6 \), which would generate terms \( <\phi_h \phi_h >_{E2N}, <\psi_h \bar{\psi}_h >_{E2N} \) to hidden matters with correct charged zero modes and other measure assignments. The nonzero intersecting number \( I_{MN} \) implies the possible existence of 1PI diagrams of multi-instantons. These effects are of higher order and will not be included in the present analysis.

4 The Higgs tree-level spectrum

In MSSM-like models, Higgs physics provides a good window to test new physics. In general, Higgs masses are sensitive to supersymmetric breaking hidden sectors. This has recently been revisited in the four-dimensional effective field theory formalism \([12]\). Earlier
discussions on this topic were present in [13]. It is shown that Higgs masses, especially the mass of $h$ can be substantially uplifted by one type of quartic couplings that were inherited from hidden sectors or extra dimensions.

In models with two Higgs doublets $H_u = (H_u^+, H_u^0)$ and $H_d = (H_d^0, H_d^-)$, there are 8 real Higgs scalars, three are eaten by the massive $W$ bosons, leaving two CP even $h$ and $H$, a CP odd $A_0$ and two charged $H^\pm$ particles. The most general form of scalar superpotential that contains operators of effective dimension less than 5 is [21, 23]

$$V = \tilde{m}_H^2 H_u^\dagger H_u + \tilde{m}_H^2 H_d^\dagger H_d - (m_{ud}^2 H_u H_d + h.c.) + \frac{\lambda_1}{2} (H_u^\dagger H_u)^2 + \frac{\lambda_2}{2} (H_d^\dagger H_d)^2$$

$$+ \lambda_3 (H_u^\dagger H_u)(H_d^\dagger H_d) + \lambda_4 (H_u^\dagger H_d)(H_d^\dagger H_u)$$

$$+ \left( \frac{\lambda_5}{2} (H_d H_d)^2 + (\lambda_6 H_u^\dagger H_u + \lambda_7 H_d^\dagger H_d)H_u H_d + h.c. \right)$$

(4.1)

where the $\mu$ term and quartic terms come from hidden sectors that break supersymmetry in the visible sector in general BMSSM models. In our model they have extra contributions of non-perturbative origin. Instead of writing the masses as functions of soft terms $\tilde{m}$, it is more convenient to express them as three new parameters. Two of them are the VEVs of $H_u^0$ and $H_d^0$, the third is $m_{A_0}$. The dimensionless parameters are in our case,

$$\lambda_1 = \lambda_2 = \frac{g'^2 + g^2}{4}, \quad \lambda_3 = \frac{g^2 - g'^2}{4},$$

$$\lambda_4 = -\frac{g^2}{2}, \quad \lambda_5 = 0, \quad \lambda_6 = \lambda_7 = 2\epsilon$$

(4.2)

The extra new parameter

$$\epsilon = \frac{A}{4!} \left( \frac{\mu}{M} \right)^2$$

(4.3)

is due to non-perturbative effects, which is in the range of $0.01 \sim 0.1$ for typical values of $M$ and $\mu$. The modifications on Higgs masses can be expressed as the the functions of $v$ and expansion of $\epsilon$.

$$\delta m_{H^\pm}^2 = 0$$

(4.4)

$$\delta m_h^2 = 2v^2 \sin(2\beta) \left( 2\epsilon + \frac{(m_{A_0}^2 + m_Z^2)\epsilon}{\sqrt{(m_{A_0}^2 - m_Z^2)^2 + 4m_{A_0}^2 m_Z^2 \sin(2\beta)}} \right) + O(\epsilon^2)$$

(4.5)

$$\delta m_H^2 = 2v^2 \sin(2\beta) \left( 2\epsilon - \frac{(m_{A_0}^2 + m_Z^2)\epsilon}{\sqrt{(m_{A_0}^2 - m_Z^2)^2 + 4m_{A_0}^2 m_Z^2 \sin(2\beta)}} \right) + O(\epsilon^2)$$

(4.6)

These operators can be constructed in five-dimensional $N = 1$ supersymmetric theory, in which the fifth dimension is compactified on the orbifold $S^1/Z_2$. The MSSM is founded to be the four-dimensional effective field theory [22].
Taking $\mu \sim 200$ GeV, $\tan\beta = 5$, the LEPII Higgs mass bound $m_h \geq 114$ GeV can be accommodated with the $\delta m_h$ at tree level when $M$ is below 20 TeV. If $\tan\beta$ decreases, one has to decrease $M$ also to uplift the $h$ mass substantially. However, for moderate value of $\tan\beta$, there will be a lower bound on $M$ from precision experiments. On the other hand, the following constraint relation between Higgs masses is unchanged,

$$m_{H_\pm}^2 = m_{A^0}^2 + m_Z^2$$

In addition to uplifting $m_h$, these operators introduce new Higgs-Higgsino interactions, which provide new channels in neutralino and chargino decays. Potentially, these phenomenological implications provides interesting tests of string theory.

5 conclusions

In this paper, we have discussed the $E_2$ instanton-induced superpotentials associated with Higgs physics in toroidal orientifolds of type IIA theories. All the moduli in flux compactifications are stabilized, which is very important to start with. Explicitly, we present a $\mathcal{N} = 1$ supersymmetric model including two $E_2$-branes. They induce the required $\mu$ term and extra quartic couplings. The later can be used to uplift the mass of the lightest Higgs boson, as expected from general analysis of four-dimensional effective field theory. In our case, they are generated by non-perturbative stringy instanton effects, instead of other mechanisms in the hidden sector.

The wrapping numbers of this model are described by two real numbers $(n_1^E, n_2^E)$, which preserve the same supersymmetry as those of $D_6$-branes. The structure of the induced quartic couplings can be calculated explicitly. For illustration, we have calculated a very simple example, in which the numbers of triangles are less than two on each torus. In this simple setting, we had the extra benefit that the mixing terms are highly suppressed. With moderate and large $\tan\beta$, the mass of the lightest Higgs boson can be uplifted substantially to meet the LEPII bound.

$^3M$ is bounded below by electro-weak precise observables (EWPO). For example, one can obtain a constraint on $M$ from the Fermi constant $G_F$, the masses of $m_W$ and $m_Z$ [24],

$$\frac{G_F}{G_F^{SM}} = \left( 1 - \left( \sin^4 \beta + 2 \sin^2 \beta - 1 \right)^{\frac{\pi^2 m_W^{(ph)^2}}{2}} \right)$$

where $\frac{G_F^{SM}}{G_F} = 1 + 0.0088 - 0.0083$ and $m_W^{(ph)} = 80.39 \pm 0.06$ GeV. For $\tan\beta = 5$, one needs $M \geq 2.17$ TeV.
Note that in all known intersecting brane models, to generate non-perturbative neutrino masses seem to forbid a non-perturbative $\mu$ term at the same time, and vice versa. Because the requirement of absence of zero modes in $E2 - E2'$ makes it very hard to satisfy all tadpole constraints and supersymmetric conditions, as point out in [5]. In our case, we have succeed to generated the $\mu$ term and extra quartic couplings, but not the desired neutrino masses. Hopefully, this can be remedied in future, without sacrificing too much other attractive features in this class of models.

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