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On fermion mass hireracies in MSSM-like quiver models with stringy corrections

Research Article

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Abstract: Using instanton effects, we discuss the problem of fermion mass hierarchies in an MSSM-like Type IIA orientifolded model with $U(3)\times Sp(1)\times U(1)\times U(1)$ gauge symmetry obtained from intersecting D6-branes. In the corresponding four-stack quiver, the different scales of the generated superpotential couplings offer a partial solution to fermion mass hierarchies. Using the known data with neutrino masses $m_{\nu_\tau} \lesssim 2$ eV, we give the magnitudes of the relevant scales.

Keywords: type IIA superstring • instanton effects and Yukawa couplings

1. Introduction

It has been recognized that important particle physics ingredients including gauge symmetry and chiral matter can be realized in Type II superstring models using the mechanism of intersecting D-branes [1, 2]. More recently, this method has progressed extensively and it is now possible to investigate some stringy effects which could give rise to deeper details seen in particle physics. In particular, D-brane instantons wrapping non trivial cycles in the internal manifold have been particularly explored in this regard [3, 4]. They give non-perturbative superpotential corrections which could explain the large mass hierarchies including the smallness of neutrino masses in the standard model (SM) [5–10].

Many works have been done along these lines using configurations realizing MSSM-like orientifolded models based on the quiver approach [11–14]. In this way, rather than considering full string theory models, working at the level of quivers allow for dealing with many important physical effects. Indeed, one could examine the possible presence or absence of couplings and other physical effects by considering quantum numbers associated with the quiver configuration data. More precisely, many investigations have been performed for MSSM-like orientifolded models satisfying necessary constraints allowing for hierarchical mass terms for all three families of quarks and leptons. This can be approached using either extra singlet fields that string theory compactifications contain or

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via Euclidean D-brane wrapping non trivial cycles producing E-instantons which we are interested in here [15–18]. They exhibit the so called uncharged modes consisting of fermion modes θ^{α}, $\overline{\theta}^{\dot{\alpha}}$ preserving half of the bulk supersymmetry breaking $N = 2$ to $N = 1$ and the four bosonic modes x^{μ}, which are associated with breakdown of four-dimensional Poincaré invariance. These kinds of uncharged zero modes appear if the instantons wrap non-rigid three-cycles in the internal manifold. However, the charged fermionic zero modes λ_a, $\bar{\lambda}_b$ appearing at the intersection between E-instantons and D-branes are crucial in the determination of superpotential corrections to four dimensional effective field theories. The non-perturbative contributions are given by performing the path integral over all fermionic zero modes. To get superpotential terms, one must ensure that all uncharged fermionic zero modes, apart from x^{μ} and θ^{α}, are projected out or lifted. In type IIA superstring for instance, this can be done by considering E2-instantons obtained from D2-branes wrapping rigid orientifold-invariant 3-cycles embedded in the internal space [3, 4]. These are referred to as rigid $O(1)$ instantons, which will be discussed here. In this picture, when instanton effects do not give rise to phenomenological inconveniences, they therefore represent a way in which one can give rise to viable mass hierarchies for the quarks and leptons as well as the smallness of neutrino masses.

In this paper we contribute to these activities by performing a discussion of mass hierarchies based on a Type IIA four-stack quiver orientifolded model in the presence of E2-instanton effects. More precisely, we focus on $\mathrm{U}(3)_a \times \mathrm{Sp}(1)_b \times \mathrm{U}(1)_c \times \mathrm{U}(1)_d$ quiver gauge theory based on intersecting D6-brane configurations with instanton corrections to the corresponding effective superpotential. For simplicity, we consider $O(1)$-instantons wrapping rigid orientifold-invariant cycles and carrying the right global $\mathrm{U}(1)_{c,d}$ charges through the $\lambda_{c,d}$, $\bar{\lambda}_{c,d}$. Then we give the corresponding MSSM-like quiver with the instanton corrections. Compared to perturbatively allowed superpotential terms of the heavy fields, the scales of the induced terms of the light fields depend on the suppression factors of E2-instantons which may induce them with some combinations with the string scale M_s through higher order terms. This mechanism offers a stringy framework to explain the fermion mass hierarchies in the quiver method. Including the left-handed neutrino masses in the analysis, we get the order of magnitude of the relevant instanton effects and recover the string scale upper bound $M_s \lesssim 10^{14}$ GeV.

The organization of this paper is as follows. In section 2, we build a quiver gauge theory realizing MSSM-like orientifolded models based on four stacks of intersecting D6-branes. In section 3, we study the fermion mass hierarchies by introducing the possible stringy corrections to the corresponding superpotential using $O(1)$ E2-instantons then we give the corresponding quiver. Their suppression magnitudes have been estimated. Section 4 will be devoted to our discussions.

2. Quiver model

In this section, we will give a quiver model based on four-stack of D6-branes embedded in IIA superstring moving on orientifolded geometry. The gauge theory lives on D6-branes wrapping four-dimensional Minkowski space and non trivial 3-cycles in the internal manifold [19, 20]. It is recalled that D6-branes carry Ramond-Ramond charges that should be canceled by the introduction of orientifold geometries related to fixed point loci of an antiholomorphic involution acting on the internal space. In such quiver models, the bifundamental chiral matter arises at the nontrivial intersection of two generic D6-branes. It turns out that, one can have symmetric or anti-symmetric tensor representations where a D6-brane intersects its image brane under the orientifold action. Furthermore, two given D6-branes might intersect in multiple points on the compact internal space, giving rise to multiple families, where the number of families is the topological intersection number of two 3-cycles belonging to middle dimensional cohomology. Indeed, a stack of N D6-branes wrapping 3-cycles gives rise to $\mathrm{U}(N)$ gauge symmetry. However, N D6-branes wrapped on cycles which are homologically-fixed or pointwise-fixed by the orientifold action give rise to $\mathrm{Sp}(N)$ and $\mathrm{SO}(N)$ gauge symmetry, respectively. Since $\mathrm{Sp}(1)$ is isomorphic to $\mathrm{SU}(2)$, the $\mathrm{SU}(2)_L$ of the MSSM-like models can be realized as a $\mathrm{Sp}(1)$ arising from a D6-brane stack wrapping an orientifold invariant 3-cycle. In this regard, the SM gauge symmetry and the matter content can be accommodated in the gauge group $\mathrm{U}(3) \times \mathrm{Sp}(1) \times \mathrm{U}(1)$ obtained from three stacks of D6-branes. Enlarging the gauge group can make the model richer. This can be done by assuming that one has an appropriate middle dimensional cohomology generated by non trivial 3-cycles. The model we consider here relies on a four stacks of D6-branes giving rise to the following gauge symmetry

$$\mathrm{U}(3)_a \times \mathrm{Sp}(1)_b \times \mathrm{U}(1)_c \times \mathrm{U}(1)_d \tag{1}$$

with a gauged flavor group $\mathrm{U}(1)_d$ distinguishing various quarks from each others. Roughly, the tadpole conditions imply the vanishing of non-abelian anomalies, while abelian and mixed anomalies are canceled via the Green-Schwarz mechanism. Generically, the anomalous $\mathrm{U}(1)^{'}s$

acquire a mass and survive only as global symmetries, which forbid various couplings on the perturbative level. Since the SM gauge symmetries contain the abelian symmetry $U(1)_Y$, and the Sp(1) does not exhibit a $U(1)_b$ which could contribute to the hypercharge, one requires that a linear combination of $U(1)_Y = \sum_k^{a,c,d} q_k U(1)_k$ remains massless. Thus, the resulting gauge group in four-dimensional spacetime can be written as

$$SU(3)_a \times Sp(1)_b \times U(1)_Y. \qquad (2)$$

Given the above gauge symmetry based on D6-brane configurations, we can construct a quiver describing MSSM-like orientifolded model. Indeed, the vanishing of anomalies, which was required to be satisfied, are used to fit the matter content involving two up quarks, one down quark charged under the $U(1)_c$ gauge symmetry and the opposite arrangement charged under $U(1)_d$. The model we present here relies on a particular intersection numbers of 3-cycles in middle dimensional cohomology. For that, we choose the following intersections[1]

$$\begin{aligned} I_{D_aD_b} = 3,\ I_{D_aD_c} = -2,\ I_{D_aD_{c^*}} = -1, \\ I_{D_aD_d} = -1,\ I_{D_aD_{d^*}} = -2,\ I_{D_dD_b} = 3, \\ I_{D_dD_{c^*}} = -3. \end{aligned} \qquad (3)$$

The other intersection numbers are set to zero. The chiral spectrum and the gauge symmetry can be represented in the table 1 together with their identification with SM matter fields. In this description, all leptons and the two Higgs doublets are charged under $U(1)_{c,d}$ symmetries while for quarks only their right handed partners are. The model with four stacks can be encoded in a quiver where each node represents a D6-brane and the links between them indicate their chiral intersections. The quiver summarizing the above spectrum with the two Higgses is shown in figure 1. Since any realistic string vacua has to exhibit the phenomenologically desired terms in four dimensional effective superpotential, we require the presence of all the MSSM Yukawa couplings and the absence of the phenomenologically undesired couplings terms on perturbative level.

In what follows, the $U(1)_{c,d}$ symmetries will be used to select the candidate terms for the superpotential. Taking into account the charges presented in the table 1, we can write down the all allowed interaction terms generating the possible Yukawa couplings. It turns out that, the only fields involved in the allowed perturbative interaction terms are those charged under the $U(1)_c$ gauge symmetry. We interpret them as the the heavy quarks c, t, b, leptons e^i and the μ term for H_u, H_d. This could be seen at the level of their associated quantum numbers

$$\begin{aligned} H_{u(0,-1,0)}q_{(1,0,0)}c_{(-1,1,0)},\ H_{u(0,-1,0)}q_{(1,0,0)}t_{(-1,1,0)}, \\ H_{d(0,1,0)}q_{(1,0,0)}b_{(-1,-1,0)},\ H_{d(0,1,0)}l^i_{(0,0,1)}e^i_{(0,-1,-1)}, \\ H_{u(0,-1,0)}H_{d(0,1,0)}. \end{aligned} \qquad (4)$$

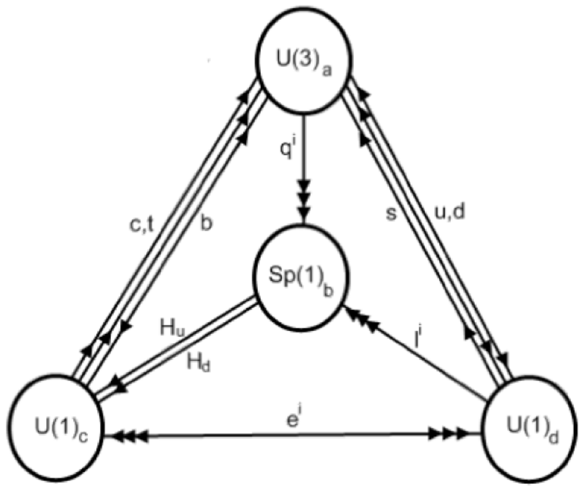

Figure 1. Quiver for intersecting D6-branes: Circles denote D6-branes while bold lines indicate the chiral spectrum. Arrows directions indicate fundamental (antifundamental) representations of U(N) gauge group.

Here the index i denotes the family index and the subscript indicates the charge under the $U(1)_{a,c,d}$ symmetries. Our quiver naturally allows for the following couplings

$$y_c H_u q\bar{c} + y_t H_u q\bar{t} + y_b H_d q\bar{b} + y_{e^i} H_d l^i \bar{e}^i + \mu H_u H_d \qquad (5)$$

where $y_{c,t,b}$ are coupling constants accounting for hierarchies between these terms. The absent couplings, which are phenomenologically desired, corresponding to fermions charged under $U(1)_d$ violate the $U(1)_{c,d}$ symmetries. These will be refereed to as the light fermions u, d, s, v that will be yielded massless, unless some kind of new effects in the defined effective field theory break the remnants abelian symmetries $U(1)_{c,d}$. In string theory, the natural candidate non-perturbative effects to violate these $U(1)_{c,d}$ symmetries are instantons arising from euclidean D-branes coupling to these fields [21–24]. They can potentially destabilize the vacuum or lead to new effects in

[1] *We have not included those involving $b^* = b$*

Table 1. The spectrum and their $U(1)_{a,c,d}$ charges for $Y=\frac{1}{6}Q_a-\frac{1}{2}Q_c-\frac{1}{2}Q_d$. The index $i=1,2,3$ denotes the family index.

$Sector$	ab	ac	ac^*	ad	ad^*	db	dc^*	bc, bc^*
$Fields$	q^i	$\overline{u}^{2,3}$	$\overline{d}^3$	$\overline{u}^1$	$\overline{d}^{1,2}$	l^i	$\overline{e}^i$	H_d, H_u
Rep	$3(3,\overline{2})$	$2(\overline{3},1)$	$1(\overline{3},1)$	$1(\overline{3},1)$	$2(\overline{3},1)$	$3(1,\overline{2})$	$3(1,1)$	$1(1,2)$
Q_a	1	−1	−1	−1	−1	0	0	0, 0
Q_c	0	1	−1	0	0	0	−1	1, −1
Q_d	0	0	0	1	−1	1	−1	0, 0
Y	1/6	−2/3	1/3	−2/3	1/3	−1/2	1	−1/2, 1/2

the four-dimensional effective action. In what follows, we will discuss the implication of such non perturbative effects in our quiver model. In particular, we will consider orientifolded invariant 3-cycles on which D2-branes can wrap to make rigid $O(1)$ instantons. Then, we discuss fermion mass hierarchies in such instanton configurations.

3. Stringy corrections and mass hierarchies

3.1. Instanton corrections

In this section we introduce the effects of spacetime instantons on the above described model. In type IIA superstring model, such non-perturbative effects are generated in principle from D0, D2 and D4-branes wrapping one, three and five-cycles respectively. However, since a Calabi-Yau manifold does not have any continuous one and five-cycles, the only relevant instantons are D2-branes wrapped on three-cycles in the internal manifold [3, 4]. Roughly, the resulting D2-instanton action takes the following form

$$S_E = S(\Lambda,\,\Phi_n) + S_E^{cl}. \tag{6}$$

In this equation, $S(\Lambda,\,\Phi_n)$ describes all terms involving fermionic instanton zero modes and four dimensional charged matter fields Φ_n while S_E^{cl} is the Dirac Born-Infeld action on D2-brane wrapping a 3-cycle Σ in the presence of the WZ term. This action reads as

$$S_E^{cl} \sim V(\Sigma) + i\int_{\Sigma} C_3 \tag{7}$$

where C_3 is the R-R 3-from coupled to the D2-brane. In four dimensions, the instanton contributions to the low energy effective theory can be obtained by performing the Grassmann path integral over all fermionic zero modes Λ:

$$S_E^{4d}(\Phi_n) = \int [D\Lambda]\, e^{-S_E} = \prod_n \Phi_n e^{-S_E^{cl}}. \tag{8}$$

The classical instanton action $e^{-S_E^{cl}}$ can absorb the U(1)'s charge excess of the matter fields operator $\prod_n \Phi_n$ [21–23]. Under such abelian symmetries the transformation property of the exponential instanton action is

$$e^{-S_E^{cl}} \longrightarrow e^{iQ(E2)\Lambda} e^{-S_E^{cl}} \tag{9}$$

where $Q(E2)$ represents the amount of the U(1)-charge violation by the $E2$-instantons. Their microscopic origin is in the extra fermionic zero modes Λ living in the intersection of the E2-branes with the $D6$ branes.

Instead of being general, let us consider a concrete configuration with $E2$-instantons wrap rigid orientifold-invariant 3-cycles. In this case, the $E2-D6$ and $E2-D6^*$ are identified. So, an $E2$-instanton intersecting $D6_{c,d}$-branes can induce the desired couplings of the matter fields operator $\prod_n \Phi_n = H_u q u$, $H_d q d$, $H_d q s$,...The $U(1)_{c,d}$ charges which are carried by the instanton action have their origin in the intersection $E2-D6_{c,d}$ pattern. For an instanton wrapping a 3-cycle with an appropriate number of times, this can be made exactly opposite to the total charge of the operator $\prod_n \Phi_n$, so the coupling (8) is also $U(1)_{c,d}$-invariant. Examining the $Q_{c,d}$ charge-excess for each term, we determine the intersection $E2-D6_{c,d}$ pattern giving rise to the right charged fermionic zero modes $\lambda_{c,d}$, $\overline{\lambda}_{c,d}$ compensating the $Q_{c,d}$ charge-excess of the desired couplings. Indeed, the up-quark coupling term

$$H_{u(0,-1,0)}\, q_{(1,0,0)}\, u_{(-1,0,1)} \tag{10}$$

violate the $U(1)_{c,d}$ symmetries by one unit as $Q_c(Hqu) = -1$, $Q_c(Hqu) = 1$. Then is generated by an instanton $E2_u$ intersecting $D6_{c,d}$ branes with the following intersection numbers

$$I_{E_u D_c} = -1,\quad I_{E_u D_d} = 1 \tag{11}$$

which gives rise to two charged modes λ_c^u, $\overline{\lambda}_d^u$:

$$Q_c(\lambda_c^u) = -N_c I_{E_u Dc} = 1,\quad Q_d(\overline{\lambda}_d^u) = -N_d I_{E_u D_d} = -1. \tag{12}$$

These charges compensate the charge-excess of the up-quark term. In this case, the corresponding E2-instanton action takes the following form

$$S_{E_u} = S^{cl}_{E_u} + \lambda_c y_u H_u q u \bar{\lambda}_d. \quad (13)$$

Performing the integration over the all fermionic zero mode, one get the desired term

$$\int d^4x\, d^2\theta\, d\lambda_c\, d\bar{\lambda}_d\, e^{-S^{cl}_{E_u} + y_u \lambda_c H_u q u \bar{\lambda}_d} = e^{-S^{cl}_{E_u}} \int d^4x\, d^2\theta\, d\lambda_c\, d\bar{\lambda}_d\, y_u \lambda_c H_u q u \bar{\lambda}_d = e^{-S^{cl}_{E_u}} y_u H_u q u. \quad (14)$$

Analogous analysis can be done for the down and strange quark coupling terms

$$H_{d(0,1,0)} q_{(1,0,0)} d_{(-1,0,-1)}, \quad H_{d(0,1,0)} q_{(1,0,0)} s_{(-1,0,-1)}. \quad (15)$$

This can be generated by an $E2_{d,s}$-instanton intersecting the $D6_{c,d}$-branes with the following intersection numbers

$$I_{E_{d,s} D_c} = 1, \quad I_{E_{d,s} D_d} = -1 \quad (16)$$

and producing the two charged modes $\bar{\lambda}^{d,s}_c$, $\lambda^{d,s}_d$. For the sector of neutrinos, the following higher order term

$$H_{u(0,-1,0)} l^i_{(0,0,1)} H_{u(0,-1,0)} l^i_{(0,0,1)} \quad (17)$$

could be generated by a $E2_\nu$-instanton intersecting $D6_{c,d}$-branes. The corresponding intersection numbers are given by

$$I_{E_\nu D_c} = -2, \quad I_{E_\nu D_d} = 2. \quad (18)$$

This leads to the charged modes λ^ν_c , $\bar{\lambda}^\nu_d$. Integrating over the all above fermionic zero modes as in (15), we get the missing superpotential terms at the four-dimensional theory

$$y_u e^{-S^{cl}_{E_u}} H_u q \bar{u} + e^{-S^{cl}_{E_{d,s}}} (y_d H_d q \bar{d} + y_s H_d q \bar{s}) + y_{\nu^i} e^{-S^{cl}_{E_{\nu^i}}} M_s^{-1} \left(H_u l^i\right)^2 \quad (19)$$

where the coupling constants $y_{d,s}$ can account for hierarchies between the d and s quarks having the same instanton suppression. The neutrinos superpotential term is highly suppressed by $1/M_s$ factor and once Higgs fields get a Vev v_u this operator gives rise directly to Majorana masses for the left-handed neutrinos depending on the scale M_s taken as the low string scale at which neutrino masses have origin.
The quiver illustrating these instanton intersection patterns with their appropriate charged fermionic zero modes is depicted in the figure 2.

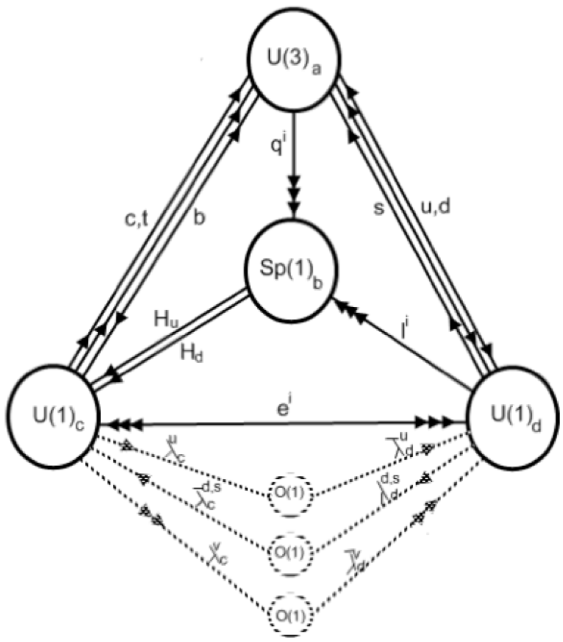

Figure 2. Quiver for intersecting D6-branes and E2-instantons: dotted circles denote the $O(1)$-instantons and dotted lines indicate their intersections with the $D6_{c,d}$-branes.

3.2. Instanton correction magnitudes

Recall that the exponential suppression effects can be derived from the classical instanton actions. Their magnitudes depend on the internal geometry on which the quiver model is based. In particular, they depend on the complex structure moduli space described by the volume of the 3-cycles wrapped by the relevant instantons. To have an idea about the induced mass hierarchies, we need to estimate the different instanton factors.Although such an approach forces consideration of deeper details of the fully defined string model, we can get an order of magnitude of the different suppression factors by refereing to the known data. Effectively, after the Higgs fields break the electroweak symmetry at the usual scale $v_u^2 + v_d^2 = (246 \text{ GeV})^2$, we take the quark couplings terms of eq. (19) together with the quark coupling terms of eq. (5) with some assumption on their scalar-fermion couplings to derive the expected magnitudes. For that, using some combination of the quark masses where their net scalar-fermion couplings effect could be neglected, one allows for getting approximate values of the corresponding suppression factors. For the u-quark suppression factor, assuming $y_u \sim y_d$ and

$y_s \sim y_c$, we find

$$e^{-S^{cl}_{E_u}} \simeq \frac{m_u m_s}{m_d m_c} \sim 3.10^{-2} \tag{20}$$

While for the d and s-quark suppression factors with $y_d + y_s \sim y_b$, we get

$$e^{-S^{cl}_{E_{d,s}}} \simeq \frac{m_d + m_s}{m_b} \sim 2.10^{-2}. \tag{21}$$

Similarly, using the known data with neutrino masses upper bound $m_{\nu\tau}$

$$\frac{e^{-S^{cl}_{E_{\nu^i}}}}{M_s} \sim 10^{-14}\ \text{GeV}^{-1}. \tag{22}$$

Considering the usual value of the string scale, namely $M_s \simeq 10^{18}$ GeV, the E2-instanton induced operator only leads to subleading corrections to the neutrino masses. Their observed order can be obtained by lowering the string scale due to large internal dimensions [25, 26]. Lowering the string scale down to the TeV scale, this leads to the interesting features at the LHC. From (22), we can deduce the string scale upper bound

$$M_s \leq 10^{14}\ \text{GeV}. \tag{23}$$

Investigating the above instanton suppression factors (20) and (22) and specializing the case that the D6-branes and the E2-branes wrap factorizable three-cycles of toroidal orientifolds, the volume of these instantons could give an insight of the detailed description of geometric background of the internal manifold. We believe that this connection deserves to be studied further.

4. Conclusion and discussion

In this paper, we have discussed the fermion mass hierarchies in a local four-stack intersecting MSSM-like D-brane quiver models using non perturbative stringy effects. More precisely, we have focused on a four-stack of D6-brane configurations in type IIA orientifolded geometries giving rise to the MSSM-like spectrum without right handed-neutrinos including rigid $O(1)$ E2-instantons. In particular, we have given a quiver model from which we have shown that some perturbatively absent coupling terms can be generated from D2-branes wrapping 3-cycles belonging to middle dimensional cohomology of the internal space.

Analyzing the quiver allows for the determination of perturbatively and non-perturbatively contributions to the superpotential using the abelian symmetries obeyed at the perturbative regime. In this approach, attributing the perturbatively realized terms to the heavy fermions and the remaining missing desired terms to the light fermions reflects interesting hierarchies. The latter are induced and exponentially suppressed by E2-instantons carrying the right charged fermionic zero modes appearing at the E2-D6 intersections.

These stringy corrections do not induce the perturbatively forbidden dangerous proton decay terms through the dimension 5 operators $q_L q_L q_L l$ and $u_R u_R d_R e_R$. Such kinds of local intersecting D-brane models have been discussed extensively in [5–7, 9]. These models, which include phenomenological constraints of fast proton decay are listed in [7]. However, the presented model does not appear in that classification since it does not involve right handed neutrinos in addition to the lepton number violation that arises through $l_L l_L e_R$ and $l_L H_u$ terms induced by the $E2_u$ instanton required to generate the missing up-quark mass term. Despite, it should be interesting to make contact with the models given in [7, 9]. We believe that this connection deserves more study. This will be reported elsewhere.

Using the known data with neutrino masses $m_{\nu_\tau} \lesssim 2$ eV, we have given the magnitudes of the relevant instanton effects and the string scale upper bound 10^{14} GeV imposed to match with the observed order of the known data.

In the end of this work, a possible discussion could be done in terms of middle dimensional cohomology of the internal space describing its complex structure. The analysis depends on the volume of the 3-cycles on which D2-branes wrap to make rigid $O(1)$ instantons. It should be very nice to find a geometrical interpretation in terms of quiver data encoded in the middle dimensional cohomology. We believe that it will be useful to explore extended Dynking geometries involving more than bosonic vertices. This seems to be promising in Type IIB D5-branes setup in the presence of D-strings wrapping 2-cycles making instantons.

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