The 750 GeV LHC diphoton excess from a baryon number conserving string model

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Abstract

We propose an explanation of the LHC data excess resonance of 750 GeV in the diphoton distribution using D-brane models, with gauged baryon number, which accommodate the Standard Model together with vector like exotics. We identify the 750 GeV scalar as either the sneutrino ($\tilde{\nu}_R$) or as an axion. Using a bottom-up approach, $\tilde{\nu}_R$ is produced via gluon fusion when scalar (supersymmetric) partners of vector like quarks and leptons are generated by demanding that the corresponding intersections respect N=1 supersymmetry. When we generate the value of $\tilde{\nu}_R$ at 750 GeV, by varying the complex structure of the torus, the string scale is limited to be in the range $10^{14} < M_s < 10^{19}$ GeV. Also, generating the Higgs scalars by imposing N=1 supersymmetry on intersections fixes naturally, to zero, the coupling of the axion to SU(2) gauge bosons $\propto F_b \wedge F_b$ and in photon-photon fusion also decouples its coupling from $G^2$ of color SU(3) by simultaneously generating the superpartners of $q_L, U_R, l_L, \nu_R$ quarks and leptons.
1 Introduction

Run 2 LHC early data from ATLAS and CMS at an energy $\sqrt{s} = 13$ TeV using integrated luminosities of 3.2 fb$^{-1}$ and 2.6 fb$^{-1}$ show hints of a new resonance in the diphoton distribution of pp collisions at an invariant mass of 750 GeV [1,2]. The corresponding excess in the cross section can be estimated to be $\sigma_{pp\to\gamma\gamma}^{13\ TeV} \sim 3-13$ fb [1,2]. At the Moriond 2016 conference, the ATLAS and CMS collaborations [3, 4, 5] updated their search, increasing the statistical significance of the excess around $m_{\gamma\gamma} \approx 750$ GeV (up to 3.9 in ATLAS and 3.4 in CMS, locally) but do not qualitatively change the main implications, still maintaining a hint for an excess at 750 GeV.

The origin of the new resonance has been attributed to a lot of different scenarios [6]-[31]. In most of the scenarios the 750 GeV resonance is a scalar (or a pseudoscalar) that is produced via gluon fusion and subsequently decays to two photons, via loops of vector-like fermions. In the context of string theories a variety of works have discussed the diphoton excess [32, 33, 34, 35, 36, 37, 38, 39, 40, 41, 42, 43, 44].

In this work, we use the $\tilde{\nu}_R$ as the source of the diphoton excess (DE) in baryon number conserving non-supersymmetric D-brane models, using the model of [45]. The presence of N=1 supersymmetry in particular intersections, gives birth to the previously massive $\tilde{\nu}_R$ that now become massless at the string scale and survives to low energies. Other works which used $\tilde{\nu}_R$ to explain DE, use an R-parity violating background in the MSSM [46, 47].

This paper is organized as follows. In Section 2 we will discuss the basic structure of the non-supersymmetric intersecting D6-brane model considered. We discuss the interpretation of the DE in terms of $\tilde{\nu}_R$ in Section 3, especially the appearance of $\tilde{\nu}_R$, and extra fermions (also scalars), in the presence of supersymmetry on intersections. In Section 4 we describe the use of the axion in the D6-brane models considered, as an alternative possibility to explain the DE. Our models possess all the ingredients to explain the use of the string axion to explain the 750 GeV state in terms of gluon or photon fusion, when superpartners of Higgsinos and some SM fermions are present because of N=1 supersymmetries preserved at particular intersections. We conclude in Section 5.
Figure 1: Quarks and leptons at the intersecting brane $U(3)_c \times U(2)_w \times U(1)_c \times U(1)_d \times U(1)_e$ model.

2 The quiver Standard Model and its embedding on a string construction

2.1 The five stack quiver Standard Model

In intersecting brane constructions chiral fermions appear as open strings stretching between brane intersecting at angles and gauge bosons living on branes. Each D-brane would give rise to a $U(1)$ and the $U(N)$ gauge group arises from $N$ overlapping D-branes (stacks). By considering a stack of D-brane configurations with $N_a, a = 1, \cdots, N$, parallel branes one gets the gauge group $U(N_1) \times U(N_2) \times \cdots \times U(N_a)$. Each $U(N_i)$ factor will give rise to an $SU(N_i)$ charged under the associated $U(1_i)$ gauge group factor that appears in the decomposition $SU(N_a) \times U(1_a)$. In this paper, we are using the five stack D6-brane string model of [45]. The initial gauge group of the model is $U(3)_c \times U(2)_b \times U(1)_c \times U(1)_d \times U(1)_e$ or $SU(3)_c \times SU(2)_w \times U(1)_b \times U(1)_c \times U(1)_d \times U(1)_e$ at the string scale. The model has the interesting properties of elevating the global symmetries of the Standard model (SM), namely Baryon B and Lepton number L, to local gauge symmetries. This model may be used to explain the diphoton excess observed from ATLAS and CMS collaborations. The representation content of the Standard Model is seen at figure [I] and table [I] charged under the five $U(1)$
symmetries $Q_a, Q_b, Q_c, Q_d, Q_e$.

| Matter Fields | Intersection | $Q_a$ | $Q_b$ | $Q_c$ | $Q_d$ | $Q_e$ | $Y$ |
|--------------|--------------|-------|-------|-------|-------|-------|-----|
| $Q_L$        | (3, 2)       | 1     | −1    | 0     | 0     | 0     | 1/6 |
| $q_L$        | 2(3, 2)      | 1     | 1     | 0     | 0     | 0     | 1/6 |
| $U_R$        | 3(3, 1)      | −1    | 0     | 1     | 0     | 0     | −2/3|
| $D_R$        | 3(3, 1)      | −1    | 0     | −1    | 0     | 0     | 1/3 |
| $L$          | 2(1, 2)      | 0     | −1    | 0     | 1     | 0     | −1/2|
| $l_L$        | (1, 2)       | 0     | −1    | 0     | 0     | 1     | −1/2|
| $N_R$        | 2(1, 1)      | 0     | 0     | 1     | −1    | 0     | 0   |
| $E_R$        | 2(1, 1)      | 0     | 0     | −1    | −1    | 0     | 1   |
| $\nu_R$      | (1, 1)       | 0     | 0     | 1     | 0     | −1    | 0   |
| $e_R$        | (1, 1)       | 0     | 0     | −1    | 0     | −1    | 1   |

Table 1: Low energy chiral fermionic spectrum of the five stack string scale $SU(3)_C \otimes SU(2)_L \otimes U(1)_a \otimes U(1)_b \otimes U(1)_c \otimes U(1)_d \otimes U(1)_e$, type I D6-brane model together with its $U(1)$ charges. Note that at low energies only the SM gauge group $SU(3) \otimes SU(2)_L \otimes U(1)_Y$ survives.

There are various gauged low energy symmetries in the models. They are defined in terms of the $U(1)$ symmetries $Q_a, Q_b, Q_c, Q_d, Q_e$, where the baryon number $B$ and lepton number $L$, respectively are equal to

$$Q_a = 3B, \quad L = Q_d + Q_e, \quad Q_a - 3Q_d - 3Q_e = 3(B - L), \quad Q_c = 2I_{3R} \quad (2.1)$$

and $I_{3R}$ being the third component of weak isospin and $3(B - L)$ and $Q_c$ are free of triangle anomalies. The $U(1)_b$ symmetry plays the role of a Peccei-Quinn symmetry, having mixed $SU(3)$ anomalies.

### 2.2 The embedding on a string construction

The interpretation of the DE in the context of string theory D-brane models, ideally, would be phenomenologically interesting if the string scale is low at the TeV region. At low scale D-brane models (LCD), extra dimensions transverse to the space where the D-branes are wrapping become large, when the string scale becomes low of order $\mathcal{O}(\text{TeV})$ and at the same time the Planck scale remains large [48, 49]. LCD’s don’t
need supersymmetry at the TeV and string constructions of low scale D5-brane models appeared in [50], [51]. Also LCD’s possess low scale string excitations, a possible signal for LHC searches [52] and also extra $Z'$ gauge bosons as the mass of $Z'$ could be at the TeV region, e.g. $M_{Z'} = (g_2^2 + g_3^2)^{1/2} v / 2 + O(v^2/M_s^2)$ [54, 55]. The present toroidal 5-stack D6-brane models don’t have large extra dimensions, but one can imagine a scenario that the six-torus can be kept small but connected to a large volume manifold.

Recently, assuming a low scale scenario, it has been shown that for a range of string scales between [10-20] TeV, the 5-stack D6-models predict the lowest $Z'$ excitation to be in the range [3.5-5.5] TeV while accommodating current anomalies in $b \rightarrow s l^+l^-$ anomalies [54].

Let us embed the SM quiver structure of table (1) in a string compactification. Our SM quiver can be embedded in a bottom-up approach in string compactification of IIA theory on a six dimensional torus equipped with an orientifolded symmetry which converts, strings with D9-branes with fluxes compactified on a six-dimensional orientifolded torus $T_6$ (where internal background gauge fluxes on the branes are turned on) by the use of a T-duality transformation on the x4, x5, x6, directions, into D6-branes intersecting at angles. In detail, we assume that the $D6_a$-branes are wrapping 1-cycles $(n_i^a, m_i^a)$ along each of the ith-$T^2$ torus of the factorized $T^6$ torus, namely $T^6 = T^2 \times T^2 \times T^2$. That means that we allow our torus to wrap factorized 3-cycles, that can unwrap into products of three 1-cycles $\Pi_a$, one for each $T^2$. We define the number of chiral fermions that are located at intersections of the branes a, b to be equal to the homology product of 3-cycles as

$$I_{ab} = [\Pi_a] \cdot [\Pi_b] = \prod_{i=1}^{3} (n_i^a m_i^b - m_i^a n_i^b)$$

and transforming in the bifundamental representation $(N_a, \bar{N}_b)$ for a left handed fermion with $I_{ab} > 0$. We also define the intersection number that determines the number of chiral fermions at the intersection of the a-brane with the orientifold image of the

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1 A light $Z'$ gauge boson could be also achieved if extra branes are added to the RR tadpoles that do not intersect with the SM branes [53].

2 We define the homology of the 3-cycle as $\Pi_a = \prod_{i=1}^{3} (n_i^a[a_i] + m_i^a[b_i])$. Because of the orientifold $\Omega R$ symmetry, where $\Omega$ is the worldvolume parity and $R$ is the reflection on the T-dualized coordinates, $T(\Omega R)T^{-1} = \Omega R$, each $D6_a$-brane 3-cycle, must have its $\Omega R$ orientifold image $\Pi_a^* = \prod_{i=1}^{3} (n_i^a[a_i] - m_i^a[b_i])$. 

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Table 2: D6-branes wrapping numbers giving rise to the standard model gauge group and chiral spectrum of table (1) at low energies. The solutions depend on five integer parameters, $n_i^a, n_i^b, n_i^c, n_i^d, n_i^e$, the NS-background $\beta^i$ and the phase parameters $\epsilon = \pm 1, \tilde{\epsilon} = \pm 1$.

| $N_i$ | $(n_i^1, m_i^1)$ | $(n_i^2, m_i^2)$ | $(n_i^3, m_i^3)$ |
|-------|------------------|------------------|------------------|
| $N_a = 3$ | $(1/\beta^1, 0)$ | $(n_a^2, \epsilon \beta^2)$ | $(3, \tilde{\epsilon}/2)$ |
| $N_b = 2$ | $(n_b^1, -\epsilon \beta^1)$ | $(1/\beta^2, 0)$ | $(\tilde{\epsilon}, 1/2)$ |
| $N_c = 1$ | $(n_c^1, \epsilon \beta^1)$ | $(1/\beta^2, 0)$ | $(0, 1)$ |
| $N_d = 1$ | $(1/\beta^1, 0)$ | $(n_d^2, 2\epsilon \beta^2)$ | $(1, -\tilde{\epsilon}/2)$ |
| $N_e = 1$ | $(1/\beta^1, 0)$ | $(n_e^2, \epsilon \beta^2)$ | $(1, -\tilde{\epsilon}/2)$ |

The fermions at the $ab^*$ intersection transform at the representation $(N_a, N_b)$ for a left handed fermion with $I_{a}b^* > 0$. Any vacuum derived from the previous intersection constraints is subject additionally to constraints coming from RR tadpole cancellation conditions [56]. That demands cancellation of D6-branes charges, wrapping on three cycles with homology $[\Pi_a]$ and O6-plane 7-form charges wrapping on 3-cycles with homology $[\Pi_O]$. Note that the RR tadpole cancellation conditions which represent the cancellation of RR charges in homology are

$$\sum_a N_a[\Pi_a] + \sum_{\alpha} N_{\alpha}[\Pi_{\alpha}] - 32[\Pi_O] = 0. \quad (2.4)$$

These conditions in string theory are stronger that the cancellation of non-abelian gauge anomalies of gauge theories, as it takes into account all the ultraviolet completion of the spectrum of the theory. The intersection numbers of the Standard Model quiver of table (1) and figure (1) are solved by the wrapping numbers seen at table (2).

Mixed U(1) gauge anomalies are cancelled via a generalized Green-Schwarz mechanism involving closed sector RR fields. Most of the five U(1)’s but two of them become massive(see eqn’s 4.3-4.6). The massless U(1)’s are

$$Q^M = n_c^1(Q_a - 3Q_d - 3Q_e) - \frac{3\tilde{\epsilon}\beta^2(n_a^2 + n_d^2 + n_e^2)}{2\beta^1}Q_c, \ (3n_a^2 + 3n_d^2 + 3n_e^2) \neq 0, \quad (2.5)$$

and

$$Q^N = \frac{3\tilde{\epsilon}\beta^2}{2\beta^1}(Q_a - 3Q_d - 3Q_e) + 19n_c^1Q_c. \quad (2.6)$$
When the condition
\[ n^1_c = \frac{\tilde{\epsilon} \beta^2}{2 \beta_1} (n^2_a + n^2_d + n^2_e), \quad n^1_c \neq 0 \]  
(2.7)
is satisfied, Q\(^M\) coincides with the Standard Model hypercharge assignment
\[ U(1)^Y = \frac{1}{6} U(1)_a - \frac{1}{2} U(1)_c - \frac{1}{2} U(1)_d - \frac{1}{2} U(1)_e. \]  
(2.8)

The tadpole conditions are
\[ \frac{9 n^2_a}{\beta_1} + 2 \frac{n^1_b}{\beta_2} + \frac{n^2_d}{\beta_1} + \frac{n^2_e}{\beta_1} + N_D \frac{2}{\beta_1^2} = 16. \]  
(2.9)

which simply adds a further constraint on the undetermined parameters of table \(2\). We have allowed the addition of \(N_D\) extra branes with wrappings \((1/\beta_1,0)(1/\beta_2,0)(2,m_D)\) that don’t intersect with the rest of the branes and thus don’t generate additional SM particles. The electroweak Higgses that are necessary for giving masses to all quarks and leptons are listed in table \(3\). Proton is stable due to the fact that baryon number

| Intersection | EW breaking Higgs | \(Q_b\) | \(Q_c\) | \(Y\) |
|--------------|------------------|--------|--------|-----|
| \(bc\)       | \(h_1\)          | 1      | -1     | 1/2 |
| \(bc\)       | \(h_2\)          | -1     | 1      | -1/2|
| \(bc^*\)     | \(H_1\)          | -1     | -1     | 1/2 |
| \(bc^*\)     | \(H_2\)          | 1      | 1      | -1/2|

Table 3: Electroweak symmetry breaking Higgses in the 5-stack D6-brane model \[45\].

\(B(\) is an unbroken gauged global symmetry surviving at low energies) anomalies cancel through a generalized Green-Schwarz mechanism(see ). In the D6-brane models with four stacks \[57\] and in the five stack models \[45\], it was noted that the Higgses arising from intersections \(bc, bc^*\) are part of the massive spectrum. They could become massless by varying the distance between the parallel branes across the 2nd tori for both intersections. In this work, we follow a different approach. The Higgs, from the NS sector, responsible for electroweak symmetry breaking will be generated by demanding that \(N=1\) supersymmetry is preserved at the intersections \(bc, bc^*\) not necessarily being the same one. Since the models are non-supersymmetric (non-SUSY) we expect that some intersections to respect some supersymmetry providing us with further constraints on the parameters \(n^1_b, n^1_c, \) etc. Related issues on \(N=1\) supersymmetries on non-SUSY
D-brane models have been discussed in \cite{61}, \cite{62}. The Yukawa interactions for the chiral spectrum of the SM's yield:
\begin{align}
Y^U_j Q_L U_R^i h_1 &+ Y^D_j Q_L D_R^i h_2 + Y^{uL}_{ij} q_L^i U_R^j h_1 + Y^{dL}_{ij} q_L^i D_R^j h_2 + \\
Y^{iL}_{hl} l_R^h \nu_R^h &+ Y^{iL}_{hl} l_R^h e_R^h h_1 + Y^{lR}_{j} L^i N_R^j h_1 + Y^{lR}_{j} L^i E_R^j h_2 + h.c
\end{align}

(2.10)

where $i = 1, 2, j = 1, 2, 3, \tilde{h} = 1$. The nature of Yukawa couplings is such that the lepton and neutrino sector of the models distinguish between different generations, e.g. the first generation from the other two generations, as one generation of neutrinos (resp. leptons) is placed on a different intersection from the other two one's. For example looking at the charged leptons of table (1) we see that one generation of charged leptons $l_L$ gets localized on $be$-intersection, while the other two generations of leptons $L$ get localized in the $bd$-intersection.

3 Interpretation of the Diphoton Excess

3.1 Di-photon preliminaries

We will consider only the case that the 750 GeV spin-0 resonance is produced from gluon decays into photons. In this case, the diphoton excess is explained via the resonant process $pp \rightarrow S \rightarrow \gamma\gamma$ where $S$ is a new uncoloured scalar boson with mass $M$, spin $J$, and width $\Gamma$ coupled to partons inside the proton. The signal cross section for the scalar mediated process mediating on shell scalar singlet $S$, is approximated as follows:

$$
\sigma(pp \rightarrow S \rightarrow \gamma\gamma) = \frac{2J+1}{sM}\Gamma(S \rightarrow gg) + \sum_q C_{qg}\Gamma(S \rightarrow q\bar{q})\Gamma(S \rightarrow \gamma\gamma).
$$

(3.1)

Assuming a spin-zero particle produced resonantly via gluon fusion, we arrive at (the last part of the following eqn. can be seen at \cite{60})

$$
\sigma(pp \rightarrow S \rightarrow \gamma\gamma) \approx K_{13} \cdot 4.92 \cdot 10^6 \text{ fb} \frac{\Gamma_{gg} \Gamma_{\gamma\gamma}}{\Gamma \Gamma} \frac{\Gamma}{M_x} \approx K_{13} \times 4.9 \cdot 10^6 \text{ fb} \frac{\Gamma_{gg} \Gamma_{\gamma\gamma}}{\Gamma \Gamma} \frac{\Gamma}{M_x}
$$

(3.2)

where we have taken into account the QCD NLO enhancement K-factors $K_{13} \approx 1.5$ \cite{13,66,67} and $\Gamma_{gg} = \Gamma(S \rightarrow gg)$, $\Gamma_{\gamma\gamma} = \Gamma(S \rightarrow \gamma\gamma)$. $\Gamma = \Gamma_{gg} + \Gamma_{\gamma\gamma}, \sqrt{s}$ are the
total width and the center of mass energy ($\sqrt{s} = 13$ TeV) respectively and $C_{gg}$ is the partonic integral

$$C_{gg} = \int_{M_{S}/s}^{1} f_g(x) f_g(\frac{M_{S}}{s x}) \frac{dx}{x},$$

where $f_g(x)$ is the function representing the gluon distribution inside the proton. The integral is computed using MSTW2008NNLO and its numerical value at 13 TeV is estimated to be $C_{gg} = 2137$. Thus in the narrow width approximation

$$\sigma(gg \rightarrow S \rightarrow \gamma\gamma) = K^{13}_{13} \times \frac{1}{M_{S} \cdot \Gamma \cdot s} C_{gg} \Gamma(S \rightarrow gg) \Gamma(S \rightarrow \gamma\gamma),$$

The partial widths $\Gamma(S \rightarrow gg), \Gamma(S \rightarrow \gamma\gamma)$ from loops involving fermions and scalars are given by (also $\Gamma(X \rightarrow gg)$)

$$\Gamma(X \rightarrow gg) = \frac{\alpha^2}{16 \pi^3} | \sum f C_{r_f} \sqrt{\tau_f y_f} \tilde{S}(\tau_f) + \sum s C_{r_s} A_{s}^2 \frac{2 M_{S}}{P(\tau_s)} |^2,$$

(3.3)

$$\Gamma(S \rightarrow \gamma\gamma) = \frac{\alpha^2}{16 \pi^3} | \sum f d_{r_f} q_{f}^2 \sqrt{\tau_f y_f} \tilde{S}(\tau_f) + \sum s d_{r_s} q_{s}^2 A_{s}^2 \frac{2 M_{S}}{P(\tau_s)} |^2,$$

(3.4)

where $a_{3} = 0.1, \alpha = 1/128, C_{r}$ is the Dynkin index of the colour representation ($C_{r} = 3$ for the triplet), $d_{r}$ is its dimension, $q_{s}$ the charge and $\sqrt{\tau_{a}} = \frac{2 m_{a}}{M_{S}},$ with $a = f, s$ for the fermion and scalar masses respectively. The functions $\tilde{S}(\tau), P(\tau)$ are

$$\tilde{S}(\tau) = 1 + (1 - \tau) f(\tau), \quad P(\tau) = \tau f(\tau) - 1,$$

$$f(\tau) = \begin{cases} \arctan^{2} \frac{1}{\sqrt{1 - \tau}} & \tau > 1 \\ -\frac{1}{4} \left( \log \frac{1 + \sqrt{1 - \tau}}{1 - \sqrt{1 - \tau}} - i \pi \right)^2 & \tau \leq 1 \end{cases}$$

(3.5)

In order to show that the D-brane model has the ingredients to explain diphoton excess, we choose the gluon fusion production cross section (3.2) at $\sigma(pp \rightarrow \gamma\gamma) \approx 3$ fb, which is the experimentally favoured value as extracted from a fit to the preferred cross sections of Morion data conference [3], [4], [5] and consider two sample cases, one of a broad resonance as favoured by ATLAS ($\Gamma = 45$ GeV) and another one of a narrow resonance as favoured by CMS. In this case, repeating the procedure of [13] and assuming that

\[^{3}\text{See relevant comment on p.29 of [66].}\]
production from $\gamma\gamma$ partons can be neglected with respect to production from $gg$, we get

\[ \text{ATLAS}, \quad \text{BR}(S \to \gamma\gamma) \cdot \text{BR}(S \to gg) \approx 6.1 \times 10^{-7} \frac{M_S}{\Gamma} \frac{\Gamma}{\Gamma_{\gamma\gamma}} = 0.06 \text{, } \Gamma = 45 \text{ GeV} \approx 1.02 \times 10^{-5} \] (3.6)

\[ \text{CMS}, \quad \text{BR}(S \to \gamma\gamma) \cdot \text{BR}(S \to gg) \approx 6.1 \times 10^{-7} \frac{M_S}{\Gamma} \frac{\Gamma}{\Gamma_{\gamma\gamma}} = 0.00013 \text{, } \Gamma = 0.1 \text{ GeV} \approx 4.69 \times 10^{-3} \] (3.7)

or equivalently

\[ \text{ATLAS}, \quad \frac{\Gamma_{gg}}{M_S} \frac{\Gamma_{\gamma\gamma}}{M_S} \approx 6.1 \cdot 10^{-7} \frac{\Gamma}{M_S} \frac{\Gamma}{\Gamma_{\gamma\gamma}} = 0.06 \text{, } \Gamma = 45 \text{ GeV} \approx 3.66 \times 10^{-8} \] (3.8)

\[ \text{CMS}, \quad \frac{\Gamma_{gg}}{M_S} \frac{\Gamma_{\gamma\gamma}}{M_S} \approx 6.1 \cdot 10^{-7} \frac{\Gamma}{M_S} \frac{\Gamma}{\Gamma_{\gamma\gamma}} = 0.00013 \text{, } \Gamma = 0.1 \text{ GeV} \approx 8.13 \times 10^{-11} \] (3.9)

where we have neglected the $K_{13}$ factor in the calculation of (3.6), (3.8), (3.7), (3.9).

### 3.2 The sneutrino $\tilde{\nu}_R$ a candidate for generating the diphoton excess

In [45] we imposed $N=1$ supersymmetry (SUSY) on the intersection $ce$ of the intersecting at angles branes $c$ and $e$, as a means to generate the spartner of $\nu_R$, namely $\tilde{\nu}_R$, which was used to break the extra U(1) (2.6) that is surviving massless beyond hypercharge at low energies. The $N=1$ SUSY condition on intersection $ce$ is

\[ ce : \quad + (\tan^{-1} \frac{\epsilon \beta_1 U_1}{n_1^1}) - (\tan^{-1} \frac{\epsilon \beta_2 U_2}{n_2^2}) - (\tan^{-1} \frac{\epsilon \beta_3 U_3}{2}) - (\frac{\pi}{2}) = 0 , \] (3.10)

which is solved (possessing the $N=1$ susy $(++)$ in the notation of [61]) by choosing

\[ \frac{\epsilon \beta_1 U_1}{n_1^1} = \frac{\epsilon U_3}{2}, \quad -\tan^{-1} \frac{\epsilon \beta_2 U_2}{n_2^2} = \frac{\pi}{2} \] (3.11)

\[ n_e^2 = 0, \quad \epsilon < 0 \] (3.12)

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\footnote{There is no $B \wedge F_i$ Chern-Simons couplings to the hypercharge $Q^M$, where $F_i$ the non-abelian field strength of the U(N), gauge bosons. As a result $Q^M \equiv Q^Y$ survives massless to low energies.}

\footnote{We define the complex structure as $U^i = \frac{m_i}{k_i}$, $i=1, 2, 3$ for the three tori. See [45] for details.}
That necessarily fixes
\[ \epsilon = -1, \; \tilde{\epsilon} = 1, \; n^1_c < 0 \] (3.13)

In this work, we interpret the observed diphoton excess as a 750 GeV scalar neutrino \( \tilde{\nu}_R \) resonance generated via gluon fusion. Since it has a low mass, it can no longer be used to break the U(1) \((2.6)\). Instead this U(1) could be broken by another gauge singlet scalar, the spartner of \( N_R \)'s located at the intersection cd which becomes massless if there is a N=1 SUSY preserved at the intersection cd (namely the ++-, N=1 susy in the notation of \([61]\)). The latter happens when the condition
\[ cd: \; + (\tan^{-1} \frac{\epsilon_1 U_1}{n^1_c}) + (\tan^{-1} \frac{2\epsilon_2 \beta_2 U_2}{n^2_d}) - (\frac{\pi}{2} + \tan^{-1} \frac{\epsilon_3 U_3}{2}) = 0 \] (3.14)
is satisfied, where we have also used \((3.11)\). Since \( \nu_R \) is located with multiplicity one at an intersection, while \( N_R \) appears with multiplicity two at another intersection, one of the spartners of \( N_R \)'s could be used to break the U(1) \((2.6)\). Also we could (as was applied recently at \([54]\)) identify \( \nu_R \) as belonging to the third generation where also we identify Q to be \((t, b)_L\) and identify q to be \((u, d)_L, (c, s)_L\); we also identify L to represent \((e, \nu_e)_L\) and \((\mu, \nu_\mu)_L\) and \( l_L \) to be \((\tau, \nu_\tau)_L\); also U, D represent \((u, c, t)_R\) and \((d, s, b)_R\) respectively. Also \((e, \nu)_R\) is identified with the right-handed electron and electron-neutrino while \((E, N)_R\) are identified with the first and second generation right-handed leptons. Thus our 5-stack D6-brane model at the string scale generates beyond the chiral spectrum of the SM and also the spartners of the neutrinos, by imposing some N=1 SUSY to be preserved at the ce intersection. There are also non-chiral fermions at the intersections \( ad, ad^*, ae, ae^*, bc, bc^*, de, be^* \) as it happens some branes to be parallel along the same \( T^2 \) torus and intersecting at angles along the rest of the tori. We discuss later those issues as matter it concerns their effect on DE.

3.3 Using fermions to saturate the decay widths \( S \to \gamma\gamma, S \to gg \)

- Contributions to \( \Gamma(S \to gg)/M_x \) from coloured vector-like fermion states

The \( \tilde{\nu}_R \) (identified as S from now on) generates the Yukawa coupling interaction\(^6\)
\[ y \; S_{(0,0,1,0,-1)} (I^x_{(-1,0,-1,0,0)}) (I^x_{(-1,0,-1,0,0)}) = y \; S^0 \; D_R^{1/3} \; X_1^{-1/3} \] (3.15)

\(^6\)The intersection numbers inside the parenthesis denotes the corresponding fermion that is localized at this intersection.
Table 4: Contributing exotic bottom quark with charge 1/3 coloured state from eqn.(3.15) to the $\Gamma_{gg}/M_x$.

| $^aX_1$ | $X_1$ Mass | $\Gamma_{gg}/M_x$ |
|--------|------------|-------------------|
| 2(3,1) | 800        | $1.7 \times 10^{-5} y^2$ |
| 2(3,1) | 900        | $1.3 \times 10^{-5} y^2$ |
| 2(3,1) | 1000       | $1.0 \times 10^{-5} y^2$ |
| 2(3,1) | 1100       | $8.8 \times 10^{-6} y^2$ |
| 2(3,1) | 1200       | $7.3 \times 10^{-6} y^2$ |
| 2(3,1) | 1500       | $4.6 \times 10^{-6} y^2$ |
| 2(3,1) | 2000       | $2.6 \times 10^{-6} y^2$ |
| 2(3,1) | 2200       | $2.1 \times 10^{-6} y^2$ |
| 2(3,1) | 2500       | $1.6 \times 10^{-6} y^2$ |
| 2(3,1) | 3000       | $1.1 \times 10^{-6} y^2$ |
| 2(3,1) | 3100       | $1.0 \times 10^{-6} y^2$ |
| 2(3,1) | 3200       | $9.9 \times 10^{-7} y^2$ |

is the only fermionic contribution to the gluon fusion partial width $\Gamma(S \rightarrow gg)$. $X_1$ is located at the intersection $ae^*$. Because $I_{ae} = 0$ at this intersection, we have generation of the non-chiral coloured states $(X_1)_{(1,0,0,0,1)} = (3,1), (X_2)_{(-1,0,0,0,-1)} = (\bar{3},1)$. The number of $X_1$’s is calculated by the non-zero intersection number in the non-parallel tori and is given by $I_{X_1} = |\beta_2(n_e^2 + n_a^2)| = |\beta_2 n_a^2|$. Choosing (for reasons clarified, later on, in the paper)

$$\beta_2 = 1, \quad n_a^2 = -2 \quad (3.16)$$

and substituting $n_e^2$ from (3.12), we have at least 2 $X_1$’s generating 2 vector pairs by pairing them with the light quarks $D_R = b_R$. Generally this mixing could influence proton decay. However, in our models baryon number is a gauge symmetry and thus proton is stable. The contribution of $b_R$ with a mass of $m_{b_R} = 4.7 \text{ MeV}$ to (3.4) is small, of the order of $10^{-10} y^2$ (assuming $y \approx O(1)$) and may not be considered further.

| Limits on the masses of vector-like bottom B quarks with electric charge -1/3 and vector couplings to W, Z, and H bosons do exist at 95 % confidence level, assuming decays |

\footnote{The superscript denotes the hypercharge assignment.}

\footnote{The other two generations of right handed down quarks, d, s provide us with small contributions as well that can be neglected.}
into standard model particles such as $B \to Wt, Zb, Hb$ at 95% C.L. Lower limits are in
the range 740-900 GeV [64] or 575-813 GeV [65] for CMS and ATLAS respectively. In
table (4) we are summarizing the contributions of $X_1$ exotic bottom quarks with masses
between 800 GeV and 3200 GeV to $\Gamma_{gg}/M_S$. We consider their Yukawa coupling a
free parameter and use $C_{X1} = 1/2$ in (3.3). We note that perturbative Type IIA
string theory allows in general for a range of $y$ including large values in the interval
$y \in [O(1), O(50)]$ [33].

- **Contributions to $\Gamma(S \to \gamma\gamma)/M_S$**

Contributions to the $S \to \gamma\gamma$ loop diagrams could come from the Yukawa mass terms
involving mixing of $S$ with the right handed fermion singlets from $cd^*$ intersection, the
$E_R$ and its vector pair from $de^*$ intersection as follows from skata

$$y_s \cdot S^{Y=0}_{(0, 0, 1, 0, -1)} \cdot (I_{cds})^{Y=+1}_{(0, 0, -1, 1, 0)} \cdot (I_{de})^{Y=-1}_{(0, 0, 0, 1, 1)} = y_s \cdot S \cdot E^{Y=1}_R \cdot (I_{de})^{Y=-1}_{(0, 0, 0, 1, 1)}$$

(3.17)

However, the number of $I_{de}$ fermions is equal to $\beta_2(2n_a^2 + n_e^2)$ [3.12, 3.14] 0. Thus there
is no contribution from (3.17) to $S \to \gamma\gamma$ width. The following vector pair of fermion
weak doublets seen in the Yukawa coupling

$$y_h S \cdot (I_{bc})^{1/2} \cdot (I_{be})^{-1/2} = y_h S \cdot (I_{be})^{1/2} \cdot l^{1/2}_L = y_h S^0 \cdot h_1^{1/2} \cdot l^{1/2}_L$$

(3.18)

contributes the to $S \to \gamma\gamma$ width. It is generated by a mixing from the higgsinos $\tilde{h}_1$
from bc intersection with the tau lepton, the lepton doublet of the third generation,
that appears with multiplicity one. The number of higgsinos is calculated from the
non-zero intersection numbers in the first and third tori

$$|I_{bc}| = |(\epsilon \beta_1)(n_b^1 + n_c^1)|, \quad n_c^1 = \frac{\beta_2}{2\beta_1} n_a^2$$

(3.19)

Making the choice

$$n_a^2 = -2, n_c^1 = -2, n_b^1 = 0, \quad \beta_1 = 1/2, \quad \beta_2 = 1 \rightarrow I_{bc} = 1$$

(3.20)

the number of higgsinos is one, thus generating one vector weak pair of doublets that
mixes with the tau lepton. Using the tau lepton mass $m_\tau = 1.78$ GeV, we find that its

---

9At the bc intersection, we have localized a non-chiral pair of Higgsinos, each of them appearing
with multiplicity given by (3.19) and (3.20). Explicitly, they are $(\tilde{h}_1)_{(0, 1, -1, 0, 0)}$, $(\tilde{h}_2)_{(0, -1, 1, 0, 0)}$.

10The choice $n_b^1 = 0$ may be justified later. See (3.24).
Table 5: Values of Yukawa couplings for the higgsino and the exotic charge 1/3 quark are obtained within the perturbative regime of string Yukawa couplings in the intersecting D6-brane model [45]. We assume $\Gamma_{\text{total}} = 0.1 \text{ GeV}$.

Then by varying the values of the colour triplet between 800-3200 GeV, we calculated the values of the higgsino couplings, such that the product $\Gamma_{gg} \Gamma_{\gamma\gamma}/M_S^2$ of eqn. (3.9) takes its LHC predicted value at $\sigma(pp \to S \to \gamma\gamma) \approx 3 \text{ fb}$ assuming a narrow width $\Gamma_{\text{total}} = 0.1 \text{ GeV}$. Table (5) summarizes our results. We find that, at $\sigma \approx 3 \text{ fb}$, with fixed higgsino mass $m_{\tilde{h}} = 800 \text{ GeV}$, only values of the exotic down quark with charge $1/3$, between [800-2000] GeV, possess perturbative string couplings (excluding the last five entries of table (5)) that are consistent with the CMS predicted diphoton excess.
Figure 2: The cross section $\sigma(pp \rightarrow X \rightarrow \gamma\gamma)$ (in fb units) in the parametric space of the Higgsino $h_1$ for a selection of the exotic down quark masses with charge 1/3 at 800, 1000, 1200 GeV. The couplings take values within perturbative string theory regime. We have set $y_s = 1$ and $y_{h_1} = 12$.

3.4 Diphoton excess in the presence of extra scalars

We have shown that a scalar superpartner of the right handed neutrino can account for the diphoton excess signal of 750 GeV and also discussed the presence of vector-like fermions which populate eqn.’s (3.3, 3.4) to explain the diphoton excess. It is also possible to generate vector-like scalar contributions to eqn.’s (3.4) by demanding that N=1 supersymmetry is preserved at the intersections $I_{ij}$ where the vector-like fermions are localized. Thus the previously massive scalars become massless at the intersections completing the N=1 chiral multiplet structure. The scalar may appear with the same multiplicity as the corresponding fermion. We will briefly discuss the conditions for these scalars to exist, without further providing further details of their contribution to $\Gamma(pp \rightarrow S \rightarrow \gamma\gamma)$.

- Contributions to $\Gamma(S \rightarrow gg)/M_S$ from coloured vector-like scalar states
Scalar contributions to $\Gamma(S \rightarrow gg)/M_S$ could be generated by demanding that in
eqn. (3.15), the intersections bc, be a N=1 supersymmetry is preserved by the corresponding branes and thus the non-chiral spartners of higgsinos the electroweak Higgses and the spartners of the the 3rd generation of leptons appear respectively. Note that as the models we discuss are not supersymmetric, these supersymmetries need not be necessarily the same. The N=1 supersymmetry condition at the intersection \( bc \) is \( \text{bc : } (-\tan^{-1} \frac{\epsilon \beta_1 U^1}{n^1_b} - \tan^{-1} \frac{\epsilon \beta_1 U^1}{n^1_c}) \pm 0 + (\tan^{-1} \frac{\bar{\epsilon} U^3_3}{2} - \frac{\pi}{2}) = 0 \), (3.23)

which is solved by eqn. (3.11) and

\[- \frac{\epsilon \beta_1 U^1}{n^1_b} = \frac{\pi}{2} \rightarrow n^1_b = 0, \quad \epsilon \in \{1, -1\} \] (3.24)

Condition (3.23) generates the electroweak scalar Higgses, superpartners of higgsinos \( \tilde{h}_1, \tilde{h}_2 \), seen at table (3).

The N=1 supersymmetry condition at the intersection be (generating the scalar superpartner of \( l_L \)) is

\[ \text{be : } \frac{\epsilon \beta_1 U^1}{n^1_b} + \frac{\epsilon \beta_2 U^2}{n^2_e} - (\frac{\tan^{-1} \frac{\epsilon U^3_3}{6}}{2} + \frac{\tan^{-1} \frac{\tilde{\epsilon} U^3_3}{2}}{2}) = 0 \], (3.25)

which is solved by

\[ \tan^{-1} \frac{\bar{\epsilon} U^3_3}{6} + \tan^{-1} \frac{\bar{\epsilon} U^3_3}{2} = \pi, \quad \bar{\epsilon} = 1 \] (3.26)

Solving (3.26) the value of complex structure moduli across the third torus is fixed at

\[ U_3 \approx 4.2 \times 10^{16} \] (3.27)

### 3.5 Estimating the mass of \( \tilde{\nu}_R \)

The mass of \( \tilde{\nu}_R \) of the yet unseen 750 GeV diphoton excess candidate in our D-brane model can be generated geometrically by varying slightly the complex structure \( U_1 \) in

\[ (\tan^{-1} \frac{\epsilon \beta_1 U^1}{n^1_b} + \tan^{-1} \frac{\epsilon \beta_1 U^1}{n^1_c}) \pm 0 - (\tan^{-1} \frac{\bar{\epsilon} U^3_3}{2} + \frac{\pi}{2}) = 0 \] (3.22)

Note that the pair of N=1 susy’s preserved at bc* intersection, namely the (++-), (+–), are different than the one’s preserved at the bc intersection.

\[ \text{Where the expressions inside the parenthesis denote the angles in the respective complex 3-dimensional orientifolded tori.} \]
Table 6: Values of complex structure variation against the string scale. We assume that the mass of $\tilde{\nu}_R$ equals 750 GeV. The preferred values of $M_s \in \{10^{14}-10^{19}\}$ GeV. Values of $\delta_1$ are compared against the value of $U_3$.

| $M_s \ (GeV)$ | $10000$ | $10^{10}$ | $10^{11}$ | $10^{12}$ | $10^{13}$ | $10^{14}$ | $10^{15}$ | $10^{16}$ | $10^{18}$ | $10^{19}$ |
|---------------|--------|--------|--------|--------|--------|--------|--------|--------|--------|--------|
| $\delta_1 \times 9.92$ | $10^{30}$ | $10^{18}$ | $10^{16}$ | $10^{14}$ | $10^{12}$ | $10^{10}$ | $10^{8}$ | $10^{6}$ | $100$ | $1$ |

the first torus. In fact, by setting its mass to 750 GeV, we will be able to set constraints on the string scale of the models. This procedure is equivalent to turning on a Fayet-Iliopoulos term in the effective theory. Assuming a slight departure of $U_1$ from its value \[ U_1 = \frac{n_1 c_1}{\beta_1} \delta_1 \] \[ \text{(3.28)} \]

where $\delta_1 \ll U_3$, we find ($M_s$ the string scale)

\[ m_{\tilde{\nu}_R}^2 = -\frac{M_s^2}{2} \frac{2\beta_1 \delta_1}{n_1^2} \] \[ \text{(3.29)} \]

As a representative example, we assume that $m_{\tilde{\nu}_R} = 750$ GeV, the string scale $M_s = 2 \times 10^{16}$ GeV and using $\beta_1 = 1/2$, $n_1^c = -2$, we find that

$\delta_1 = 2.48 \times 10^6$ \[ \text{(3.30)} \]

is really a small number compared to $U_3$ value \[ \text{(3.27)} \]. Eqn. \[ \text{(3.28)} \] also fixes moduli $U_1$. Varying the string scale from $7 \text{ TeV} < M_s < 10^{19}$ GeV, we list at table (6) that small numbers (less or equal to 1% of $U_3$) for the variation $\delta_1$ are produced only for $10^{14} < M_s < 10^{19}$ GeV and thus a high string scale is preferred. For the parameter values we have used, namely $n_1^2 = n_2^2 = n_1^1 = 0$, $-\epsilon = \tilde{\epsilon} = 1$, $\beta_1 = 1/2$, $\beta_2 = 1$ the tadpole condition \[ \text{(2.9)} \] is satisfied for $N_D = 13$ hidden anti-branes. Thus the wrappings of the extra branes become $(1/\beta_1, 0)(1/\beta_2, 0)(2, -m_D)$. The high string scale result was expected as at the present string models have a naturally high string scale since there are no transverse dimensions to the branes that can be made large by lowering the string scale at the TeV \[ \text{[48], [49].} \]
4 The axion as a string theory 750 GeV candidate responsible for di-photon excess

In \cite{41, 34, 35, 36} the axion was discussed as a solution to the diphoton excess problem in the context of string theory. The parameters of our classes of D6-brane string models, accommodate these scenaria when certain superpartners of SM particles are present.

The effective axion Langrangian for an axion $\Phi$ coupled to gluons and photons for our toroidal models is (see \cite{36} for relevant discussion)

$$\mathcal{L}_{a0} = \frac{\alpha_s}{4\pi} g g_s^{-1} G_{\mu\nu} \tilde{G}^{\mu\nu} + \frac{\alpha_Y}{4\pi} g g_Y^{-1} B_{\mu\nu} \tilde{B}^{\mu\nu},$$

(4.1)

where $f$ the axion decay coupling constant; $\alpha_s, \alpha_Y$ the strong and hypercharge fine structure constants; $g_s, g_Y$ are model dependent coefficients. In the context of the toroidal models we are discussing there are four potential candidate axions $\Phi_i, i = 0, 1, 2, 3$. Their duals, the RR sector two forms $B_{i2}$ (and their four dimensional Poincare dual axion scalars $\Phi_i$) couple to all the U(1) field strengths $F_i$ \cite{57} (part of the original U(N) gauge groups of the different brane stacks) as

$$\sum_\alpha k^i_\alpha (B_i \wedge \text{tr}(F^\alpha)), \quad (4.2)$$

In the present models the five different U(1)'s couple as \cite{45}:

$$B_2^1 \wedge \left( \frac{-2\epsilon\beta^1}{\beta^2} \right) F^b,$$

$$B_2^2 \wedge \left( \frac{\epsilon\beta^2}{\beta^1} \right) (9F^a + 2F^d + F^e),$$

$$B_2^3 \wedge \left( \frac{3\epsilon n_a}{2\beta^1} F^a + \frac{n_b}{\beta^2} F^b + \frac{n_c}{\beta^2} F^c - \frac{\epsilon n_d}{2\beta^1} F^d - \frac{\epsilon n_e}{2\beta^1} F^e \right).$$

(4.3)

Notice that $B \wedge F$ couplings induce a Stueckelberg mass term for the anomalous U(1)'s that has a non-zero $k^i_\alpha$. The three U(1)'s that couple to $B_2^i, i=1,2,3$, cancel their triangle anomalies, receive a mass and the corresponding U(1)'s remain as global symmetries to low energies. The fourth U(1) combination, the hypercharge \cite{2.5, 2.7, 2.9} also remains massless as it does not couple to any $F^i$'s. The coupling of $B_2^0$ to any $F^i$ is zero (as we have imposed the condition $\Pi_{i=1}^3 m^i = 0$ to all branes by construction to this class of models) and thus the associated axion $\Phi^0$ that stays massless. The associated axions
\( \Phi^i \) from the closed string sector couple to the U(N) field strengths as

\[
\sum_{\alpha} \lambda^i_{\alpha} \Phi^i \text{tr}(F^\alpha \wedge F^\alpha), \quad \alpha = a, b, c, d, e; \quad i = 0, 1, 2, 3
\]  \hspace{1cm} (4.4)

cancelling the mixed U(1) triangle anomalies \( A^{ij} \) of the massive U(1)’s to the gauge groups as \( A^{ij} + k^i_{\alpha} \lambda^j_{\alpha} = 0 \). In fact, the lagrangian coupling of the axion becomes

\[
\Phi_o \left( \lambda^a_{o} F_a \wedge F_a + \lambda^b_{o} F_b \wedge F_b + \lambda^d_{o} F_d \wedge F_d + \lambda^e_{o} F_e \wedge F_e \right),
\]  \hspace{1cm} (4.5)

where \( \lambda^a_{o} \) are model dependent coefficients. In our models, the non-zero axion-like couplings are

\[
\Phi^1 \wedge \left[ \frac{\epsilon \epsilon \beta^2}{2 \beta^1} (F^a \wedge F^a) - \frac{\epsilon \epsilon \beta^2}{\beta^1} (F^d \wedge F^d) - \frac{\epsilon \epsilon \beta^2}{2 \beta^1} F^e \wedge F^e \right],
\]

\[
\Phi^2 \wedge \left[ \frac{-\epsilon \beta^1}{2 \beta^2} (F^b \wedge F^b) + \frac{\epsilon \beta^1}{\beta^2} (F^e \wedge F^e) \right],
\]

\[
\Phi^o \wedge \left( \frac{3n^2_{a}}{\beta^1} (F^a \wedge F^a) + \frac{\epsilon n^1_{b}}{\beta^2} (F^b \wedge F^b) + \frac{n^2_{d}}{\beta^1} (F^d \wedge F^d) + \frac{n^2_{e}}{\beta^1} (F^e \wedge F^e) \right),
\]  \hspace{1cm} (4.6)

The present models offer a variety of possibilities as matter as it concerns the possible couplings of axion to QCD field strength and the photon. We list them as follows:

### 4.1 Producing the axion via photon fusion

In section \([3]\), sneutrino was the 750 GeV candidate that was produced via gluon fusion. Assuming at this section that the axion (and not the \( \tilde{\nu}_R \)) could be the 750 GeV candidate, the scalar \( \tilde{\nu}_R \) could be used to break the extra beyond hypercharge U(1) \((2.6)\). Let us further assume that the axion is produced via photon fusion. This possibility has been discussed in the four stack models of \([34]\) assuming a low scale string theory. In our 5-stack parametric classes of string models, photon fusion is easily accommodated if the coupling of the axions to the SU(2), SU(3) field strengths become zero, namely

\[
n^2_a = 0, \quad n^1_b = 0, \quad n^2_e = 0
\]

\( \text{(or} \quad n^2_a = 0, \quad n^1_b = 0, \quad n^2_d = 0) \)  \hspace{1cm} (4.7)

We have seen in \([3.24]\) that the coefficient \( n^1_b \), which describes the axion coupling to the SU(2) gauge bosons, can be relaxed to zero as a result of imposing N=1 supersymmetry on the intersections \( bc, bc* \), thus generating the Higgs necessary to give masses to all
quarks and fermions. The condition $n^2_e = 0$ is derived in (3.12) by demanding that a N=1 supersymmetry is preserved at the intersection $ce$, in order to generate a massless sneutrino $\tilde{\nu}_1$ that is used to break the $U(1)$ (2.6). Finally, the condition $n^2_a = 0$ is necessary if we want the axion to be produced by photon fusion. In this case (as in [34], [35]), the axion should not couple to colour $SU(3)$ gauge bosons to avoid unwanted diphoton signals. We derive this condition by demanding that the spartner of right handed up quark $U_R$ is generated when N=1 supersymmetry is preserved at the intersection $ac$ where the spartner of $U_R$ is localized. The supersymmetry condition at $ac$ is

$$ac: \left( \tan^{-1} \frac{\epsilon \beta_1 U_1}{n^2_c} \right) + \left( \tan^{-1} \frac{\epsilon \beta_2 U_2}{n^2_a} \right) + \left( \tan^{-1} \frac{\tilde{\epsilon} U_3}{6} - \frac{\pi}{2} \right) = 0 \quad (4.8)$$

The branes $a, c$ preserve the N=1 supersymmetry (-++). Eqn. (4.8) is solved by (3.11), (3.13) and

$$\tan^{-1} \frac{\epsilon \beta_2 U_2}{n^2_a} = \frac{\pi}{2} - \tan^{-1} \frac{\tilde{\epsilon} U_3}{6} - \tan^{-1} \frac{\tilde{\epsilon} U_3}{2} \quad (4.9)$$

Eqn. (4.9) is solved by

$$n^2_a = 0 \quad (4.10)$$

Using also that $\epsilon$ is negative (see eqn. (3.24)) results in the condition (3.26), that was derived by demanding that the spartner of $l_L$ is generated at the intersection $be$, fixing the $U_3$ modulus at its value (3.27). Let us assume that the same supersymmetry (-++) is preserved at the $ab^*$ intersection. Then conditions (4.10) and also (3.26) which solve (4.10) at intersection $ac$, also solve the N=1 supersymmetry condition at $ab^*$, namely

$$ab^*: \left( \tan^{-1} \frac{\epsilon \beta_1 U_1}{n^2_b} \right) + \left( \tan^{-1} \frac{\epsilon \beta_2 U_2}{n^2_a} \right) + \left( \tan^{-1} \frac{\tilde{\epsilon} U_3}{6} + \tan^{-1} \frac{\tilde{\epsilon} U_3}{2} \right) = 0 \quad (4.11)$$

Condition (4.11) generates the spartner of $q_L$.

### 4.2 Producing the axion via gluon fusion

Axionic gluon fusion (AGF) has been discussed in a string theory context in [41], [36]. The axion AGF could be present in our models if a non-zero coupling of axion to the $SU(3)^2$ field strength and the photon $F^2$ exists. Thus we need

$$n^1_b = 0, \quad n^2_e = 0 \quad (or \quad n^1_b = 0, \quad n^2_e = 0) \quad (4.12)$$

Alternatively, we could have let $n^2_e \neq 0$, in which case, necessarily, since we have to break the extra $U(1)$ (2.6), we will generate the required gauge singlet scalars by enforcing N=1 supersymmetry at the de intersection (see eqn. (3.14)). This procedure generates the two sneutrinos $\tilde{N}_R$, superpartners of $N_R$, both localized at the same intersection. One of them, could be used to break the $U(1)$ (2.6).
In [36] the value of \( n_{b} \) was allowed from RR tadpole conditions to be chosen equal to zero. In our models the conditions \([4.12]\) are obtained, as we require the presence of Higgses, higgsinos and \( \tilde{\nu}_R \) simultaneously with the SM quarks and leptons. As we have already discussed, this is achieved with the simultaneous presence of supersymmetry on intersections, bc, bc* and ce respectively. As seen from \([4.1], [4.6]\), the QCD coefficient \( g_s \) in \([4.1]\) is proportional to the \( n_a^2 \) coefficient. Similarly, the hypercharge related \( g_Y \) in \([4.1]\) is related to the \( n_a^2, n_d^2, n_e^2 \) coefficients since the hypercharge depends on \( U(1)^a, U(1)^d, U(1)^e \). Values of the ratios \( g_s/f, g_Y/f \) constrained by dijet constraints, such that gluonic fusion is achieved where further discussed in [36].

5 Conclusions

We have shown that the sneutrino can explain the 750 GeV diphoton excess, produced via gluon fusion, within the context of high scale D-brane string models. Bottom quarks with charges 1/3 contribute to the diphoton excess and mix with the SM down quarks together with higgsinos. We have not considered the contribution of scalars to the diphoton excess in detail in this work. The mass of the sneutrino can be naturally as low as 750 GeV, as a result of variation of the complex structure modulus \( U_3 \) on a 2-dimensional torus (turing on a Fayet Iliopoulos term). The variation is small (at least 1 % of \( U_3 \)) only for values of the string scale \( M_s \in [10^{14}, 10^{19}] \) GeV. We have also discussed the possibility that the axion is responsible for the diphoton excess. In the present class of models an axion always remain massless as it does not get a mass from the generalized Green-Schwarz mechanism involving the Stuelkeberg BF couplings. In this case, the diphoton signal could be produced either from photon fusion or gluon fusion. The couplings of the axions to SU(2) gauge bosons becomes naturally zero, as a result of the supersymmetry condition involving the angles between the branes b, c and b, c*, that generates the scalar Higgses needed for electroweak symmetry breaking. We have also shown that in the case where no coupling between the axions and the colour SU(3) \( G^2 \) (as considered in [35]) exists, the zero interaction result can be justified in our non-supersymmetric D6-brane model, due to the presence of N=1 supersymmetry in the intersections \( ab*, ac, be, ce \), where \( q_L, U_R, l_L, \nu_R \) are localized respectively. Thus the superpartners of the \( q_L, U_R, l_L, \nu_R \) should also be present in the low energy effective action of the 5-stack D6-brane model.
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