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The 750 GeV Diphoton Excesses in a Realistic D-brane Model

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Abstract

We study the diphoton excesses near 750 GeV recently reported by the ATLAS and CMS collaborations within the context of a phenomenologically interesting intersecting/magnetized D-brane model on a toroidal orientifold. It is shown that the model contains a SM singlet scalar as well as vector-like quarks and leptons. In addition, it is shown that the singlet scalar has Yukawa couplings with vector-like quarks and leptons such that it may be produced in proton-proton collisions via gluon fusion as well as decay to diphotons through loops involving the vector-like quarks. Moreover, the required vector-like quarks and leptons may appear in complete $SU(5)$ multiplets so that gauge coupling unification may be maintained. Finally, it is shown that the diphoton signal may be accommodated within the model.

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I. INTRODUCTION

Recently, the ATLAS [1] and CMS [2] collaborations have both reported an excess in the diphoton channel near 750 GeV. With an integrated luminosity of 3.2 fb^{-1}, the ATLAS collaboration has observed a local 3.6σ excess at a diphoton invariant mass of around 747 GeV, assuming a narrow width resonance. For a wider width resonance, the signal significance increases to 3.9σ with a preferred width of about 45 GeV. With an integrated luminosity of 2.6 fb^{-1}, the CMS collaboration has also observed a diphoton excess with a local significance of 2.6σ at invariant mass of around 760 GeV. Assuming a decay width of around 45 GeV, the significance reduces to 2σ in this case. The corresponding excesses in the cross section can be roughly estimated as $\sigma^{13\ \text{TeV}}_{pp\to\gamma\gamma} \sim 3-13$ fb [1, 2]. While this is well below the threshold to claim a discovery, this excess could be the first signal of physics beyond the Standard Model. As such, it is worthwhile to consider possible models of new physics which may explain the excess. Indeed, many groups have proposed such possible explanations [3–23]

Perhaps the simplest explanation for the excess is the addition of a SM singlet scalar with a mass near 750 GeV along with additional vector-like multiplets of colored particles. With this set-up, the singlet may be produced via gluon fusion with the vector-like particles appearing in loops. Similarly, the singlet may decay to diphotons. However, in order to preserve gauge coupling unification as in supersymmetric versions of the SM, these vector-like states should come in complete multiplets of $SU(5)$. Moreover, to preserve unification and avoid Landau poles, the types and numbers of $SU(5)$ multiplets is restricted [14].

Such light vector-like multiplets are often found in models constructed within the framework of string theory[24]. Indeed, vector-like states are generically present in intersecting/magnetized D-brane models on orientifold backgrounds [25–37]. One such model satisfying all global consistency conditions has been constructed from intersecting/magnetized D-branes within the context of Type II orientifold compactifications [38, 39] on a $T^6/(\mathbb{Z}_2\times\mathbb{Z}_2)$ background. This model corresponds to the MSSM with three generations of quarks and leptons as well as a single pair of Higgs fields. The model contains a minimal amount of exotic matter, which may be decoupled from the low-energy sector. In addition, the tree-level gauge couplings are automatically unified at the string scale [38, 39]. Finally, the Yukawa couplings are allowed by global $U(1)$ symmetries, and it is possible to obtain correct masses and mixings for quarks and charged leptons. Thus, this is a phenomenologically interesting model worthy of detailed study.

In the following, we briefly summarize the intersecting D-brane model under study, which is a variation of the model discussed above. It is shown that vector-like quarks are present in the model, and that these states may appear in complete multiplets of $SU(5)$ so that gauge coupling unification may be maintained. Furthermore, it is shown that there are SM singlets in the model which have Yukawa couplings to the vector-like quarks and leptons. It should be emphasized that this is a non-trivial result as these fields carry global $U(1)$ charges, under which the Yukawa coupling must be neutral. Finally, we show that the diphoton excesses can be accomodated in the model.

II. THE MODEL

A phenomenologically interesting intersecting D-brane model has been studied in Refs. [38, 39]. A variation of this model with a different hidden sector was also studied in Refs. [40, 41]. Type IIA orientifold string compactifications with intersecting D-branes (and their Type IIB duals with magnetized D-branes) have provided exciting geometric tools with which the MSSM may be engineered. While this approach may not allow a first-principles understanding of why the SM gauge groups and associated matter content arises, it may allow a deeper insight into how the finer phenomenological details of the SM may emerge. In short, D6-branes in Type IIA fill (3+1)-dimensional Minkowski spacetime and wrap 3-cycles in the compactified manifold, such that a stack of N branes generates a gauge group $\mathrm{U}(N)$ [or $\mathrm{U}(N/2)$ in the case of $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$] in its world volume.

In general, the 3-cycles wrapped by the stacks of D6-branes intersect multiple times in the internal space, resulting in a chiral fermion in the bifundamental representation localized at the intersection between different stacks a and b. The multiplicity of such fermions is then given by the number of times the 3-cycles intersect. Each stack of D6-branes a may intersect the orientifold images of other stacks b', also resulting in fermions in bifundamental representations. Each stack may also intersect its own image a', resulting in chiral fermions in the symmetric and antisymmetric representations. Non-chiral matter may also be present between stacks of D-branes which do not intersect on one two-torus. A zero intersection number between two stacks of branes implies that the branes are parallel on at least one torus. At such kind of intersection additional non-chiral (vector-like) multiplet pairs from $ab + ba$, $ab' + b'a$, and $aa' + a'a$ can arise. Global consistency of the model requires certain constraints to be satisfied, namely, Ramond-Ramond (R-R) tadpole cancellation and the preservation of $\mathcal{N} = 1$ supersymmetry. In particular, the conditions for preserving $\mathcal{N} = 1$ supersymmetry fixes the complex structure parameters.

The set of D6 branes wrapping the cycles on a $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold shown in Table I results in a three-generation Pati-Salam model with additional hidden sectors. The full gauge symmetry of the model is given by $[\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{U}(2)_R]_{\text{observable}} \times [\mathrm{U}(2) \times \mathrm{USp}(2)^2]_{\text{hidden}}$. As discussed in detail in [38, 39], with this configuration of D6 branes all R-R tadpoles are canceled, K-theory constraints are satisfied, and $\mathcal{N} = 1$ supersymmetry is preserved. Furthermore, the tree-level MSSM gauge couplings are unified at the string scale. Finally,

TABLE I. D6-brane configurations and intersection numbers for a three-family Pati-Salam model on a Type-IIA $T^6/(\mathbb{Z}_2\times\mathbb{Z}_2)$ orientifold, with a tilted third two-torus. The complete gauge symmetry is $[\mathrm{U}(4)_C\times\mathrm{U}(2)_L\times\mathrm{U}(2)_R]_{\text{observable}}\times[\mathrm{U}(2)\times\mathrm{USp}(2)^2]_{\text{hidden}}$ and $\mathcal{N}=1$ supersymmetry is preserved for $\chi_1=3$, $\chi_2=1$, $\chi_3=2$.

	$\mathrm{U}(4)_C\times\mathrm{U}(2)_L\times\mathrm{U}(2)_R\times\mathrm{U}(2)\times\mathrm{USp}(2)^2$											
	N	$(n^1,l^1)\times(n^2,l^2)\times(n^3,l^3)$	n_S	n_A	b	b'	c	c'	d	d'	3	4
a	8	$(0,-1)\times(1,1)\times(1,1)$	0	0	3	0	-3	0	0(2)	0(1)	0	0
b	4	$(3,1)\times(1,\ 0)\times(1,-1)$	2	-2	-	-	0(6)	0(1)	1	0(1)	0	-3
c	4	$(3,-1)\times(0,1)\times(1,-1)$	-2	2	-	-	-	-	-1	0(1)	3	0
d	4	$(1,0)\times(1,-1)\times(1,1)$	0	0	-	-	-	-	-	-	-1	1
3	2	$(0,-1)\times(\ 1,\ 0)\times(\ 0,\ 2)$	$\chi_1=3$									
4	2	$(0,-1)\times(\ 0,\ 1)\times(\ 2,0)$	$\chi_2=1,\ \chi_3=2$									

the Yukawa matrices for quarks and leptons are rank 3 and it is possible to obtain correct mass hierarchies and mixings.

Since $\mathrm{U}(N)=\mathrm{SU}(N)\times\mathrm{U}(1)$, associated with each the stacks a, b, c, and d are U(1) gauge groups, denoted as $\mathrm{U}(1)_a$, $\mathrm{U}(1)_b$, $\mathrm{U}(1)_c$, and $\mathrm{U}(1)_d$. In general, these U(1)s are anomalous. The anomalies associated with these U(1)s are canceled by a generalized Green-Schwarz (G-S) mechanism that involves untwisted R-R forms. As a result, the gauge bosons of these Abelian groups generically become massive. However, these U(1)s remain as global symmetries to all orders in perturbation theory. These global U(1) symmetries may also result in the forbidding of certain superpotential operators, such as Yukawa couplings and those which mediate baryon and lepton number violation. However, these *global* symmetries may be broken by nonperturbative effects, such as from D-brane instantons.

Some linear combinations of $U(1)$s may also remain massless if certain conditions are satisfied. For the present model, precisely one linear combination has a massless gauge boson and is anomaly- free:

$$\mathrm{U}(1)_X=\mathrm{U}(1)_a+2\left[\mathrm{U}(1)_b+\mathrm{U}(1)_c+3\mathrm{U}(1)_d\right]. \tag{1}$$

Thus, the effective gauge symmetry of the model at the string scale is given by

$$\mathrm{SU}(4)_C\times\mathrm{SU}(2)_L\times\mathrm{SU}(2)_R\times\mathrm{U}(1)_X\times\left[\mathrm{SU}(2)\times\mathrm{USp}(2)^2\right]. \tag{2}$$

The gauge symmetry is first broken by splitting the D-branes as $a\to a1+a2$ with $N_{a1}=6$ and $N_{a2}=2$, and $c\to c1+c2$ with $N_{c1}=2$ and $N_{c2}=2$, and $d\to d1+d2$ with$N_{d1}=2$

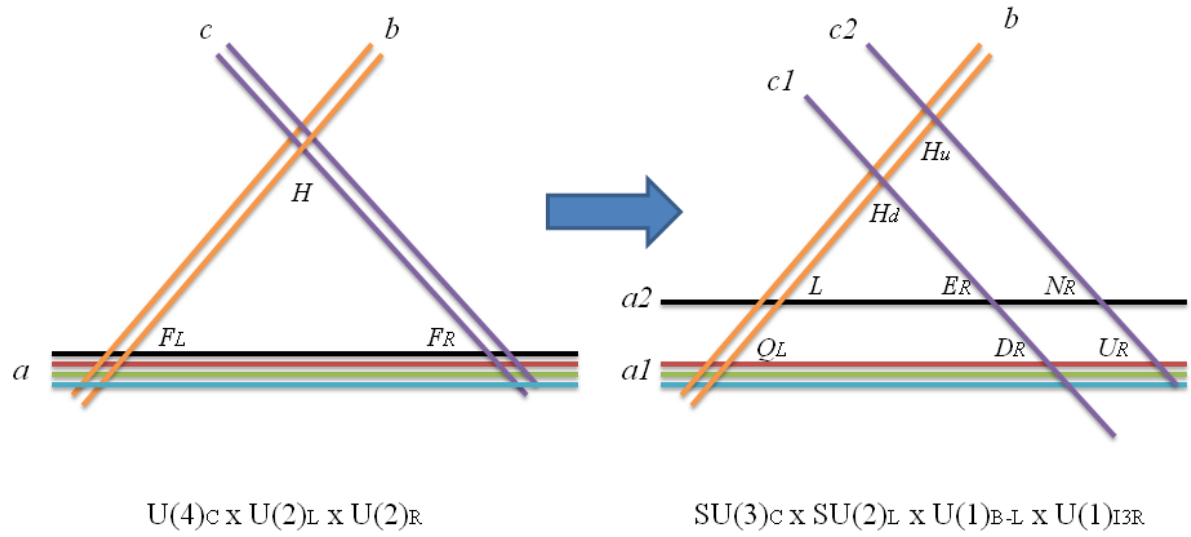

FIG. 1. Breaking of the effective gauge symmetry via D-brane splitting. This process corresponds to assigning VEVs to adjoint scalars, which arise as open-string moduli associated with the positions of stacks a and c in the internal space.

TABLE II. The chiral superfields, their multiplicities and quantum numbers under the gauge symmetry $[\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y]_{\text{observable}} \times [\mathrm{U}(1)_{I3V} \times \mathrm{USp}(2)^2]_{\text{hidden}}$.

	Mult.	Quantum Number	Q_{I3R}	Q_{B-L}	Q_{3B+L}	Q_Y	Q_B	Q_L	$Q_{Y'}$	Field
$a1b$	3	$(3, \overline{2}, 1, 1, 1, 1)$	0	1/3	1	1/6	1/3	0	1/6	Q_L
$a1c2$	3	$(\overline{3}, 1, 1, 1, 1, 1)$	-1/2	-1/3	-1	-2/3	-1/3	0	-2/3	U_R
$a1c1$	3	$(\overline{3}, 1, 1, 1, 1, 1)$	1/2	-1/3	-1	1/3	-1/3	0	1/3	D_R
$a2b$	3	$(1, \overline{2}, 1, 1, 1, 1)$	0	-1	1	-1/2	0	1	-1/2	L
$a2c1$	3	$(1, 2, 1, 1, 1, 1)$	1/2	1	-1	1	0	-1	1	E_R
$a2c2$	3	$(1, 2, 1, 1, 1, 1)$	-1/2	1	-1	0	0	-1	0	N_R

and $N_{d2} = 2$, as shown schematically in Fig.1. After splitting the D6-branes, the gauge symmetry of the observable sector is

$$SU(3)_C \times SU(2)_L \times U(1)_{I3R} \times U(1)_{B-L} \times U(1)_{3B+L} \times U(1)_{I3V}, \tag{3}$$

TABLE III. The chiral hidden sector superfields, their multiplicities and quantum numbers under the gauge symmetry $[\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y \times]_{\text{observable}} \times [\mathrm{U}(1)_{I3V} \times \mathrm{USp}(2)^2]_{\text{hidden}}$.

	Mult.	Quantum Number	Q_{I3R}	Q_{B-L}	Q_{3B+L}	Q_Y	Q_{I3V}	Q_B	Q_L	$Q_{Y'}$	Field
$bd1$	1	$(1, \overline{2}, 1, 1, 1, 1)$	0	0	-4	0	1/2	-1	-1	1/2	X_{bd1}
$bd2$	1	$(1, \overline{2}, 1, 1, 1, 1)$	0	0	-4	0	-1/2	-1	-1	-1/2	X_{bd2}
$c1d1$	1	$(1, 1, 1, 1, 1, 1)$	1/2	0	4	1/2	1/2	1	1	1	X_{c1d1}
$c1d2$	1	$(1, 1, 1, 1, 1, 1)$	1/2	0	4	1/2	-1/2	1	1	0	X_{c1d2}
$c2d1$	1	$(1, 1, 1, 1, 1, 1)$	-1/2	0	4	-1/2	1/2	1	1	0	X_{c2d1}
$c2d2$	1	$(1, 1, 1, 1, 1, 1)$	-1/2	0	4	-1/2	-1/2	1	1	-1	X_{c2d2}
$b4$	3	$(1, \overline{2}, 1, 1, 1, 2)$	0	0	-2	0	0	-1/2	-1/2	0	X^i_{b4}
$c13$	3	$(1, 1, \overline{2}, 1, 2, 1)$	1/2	0	-2	1/2	0	-1/2	-1/2	1/2	X^i_{c13}
$c23$	3	$(1, 1, \overline{2}, 1, 2, 1)$	-1/2	0	-2	-1/2	0	-1/2	-1/2	-1/2	X^i_{c23}
$d13$	1	$(1, 1, 1, 1, 2, 1)$	0	0	-6	0	1/2	-3/2	-3/2	1/2	X_{d13}
$d23$	1	$(1, 1, 1, 1, 2, 1)$	0	0	-6	0	-1/2	-3/2	-3/2	-1/2	X_{d23}
$d14$	1	$(1, 1, 1, 1, \overline{2})$	0	0	-6	0	1/2	-3/2	-3/2	1/2	X_{d14}
$d24$	1	$(1, 1, 1, 1, \overline{2})$	0	0	-6	0	-1/2	-3/2	-3/2	-1/2	X_{d24}
b_S	2	$(1, 3, 1, 1, 1, 1)$	0	0	-4	0	0	-1	-1	0	T^i_L
b_A	2	$(1, 1, 1, 1, 1, 1)$	0	0	4	0	0	1	1	0	S^i_L
c_S	2	$(1, 1, 1, 1, 1, 1)$	0	0	-4	0	0	0	-1	0	T^i_R

where

$$U(1)_{I3R} = \frac{1}{2}(U(1)_{c1} - U(1)_{c2}), \quad U(1)_{B-L} = \frac{1}{3}(U(1)_{a1} - 3U(1)_{a2}),$$
$$U(1)_{I3V} = \frac{1}{2}(U(1)_{d1} - U(1)_{d2}), \tag{4}$$

and

$$U(1)_{3B+L} = -[U(1)_{a1} + U(1)_{a2} + 2(U(1)_b + U(1)_{c1} + U(1)_{c2} + 3U(1)_{d1} + 3U(2)_{d2})], \tag{5}$$

The gauge symmetry must be further broken to the SM, with the possibility of one or more additional U(1) gauge symmetries. In particular, the $U(1)_{B-L} \times U(1)_{I_{3R}} \times U(1)_{3B+L}$ gauge symmetry may be broken by assigning VEVs to the right-handed neutrino fields N^i_R. In this case, the gauge symmetry is broken to

$$[SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_B]_{observable} \times \left[U(1)_{I3V} \times USp(2)^2\right]_{hidden} \tag{6}$$

TABLE IV. The vector-like superfields not charged under $U(1)_{I3V}$, their multiplicities and quantum numbers under the gauge symmetry $[SU(3)_C \times SU(2)_L \times U(1)_Y]_{observable} \left[U(1)_{I3V} \times USp(2)^2\right]_{hidden}$ and their charges under different $U(1)$ groups. Here a, b, c, etc. refer to different stacks of D-branes.

	Mult.	Quantum Number	Q_{I3R}	Q_{B-L}	Q_{3B+L}	Q_Y	Q_B	Q_L	$Q_{Y'}$	Field
$a1b'$	3	$(3, 2, 1, 1, 1, 1)$	0	1/3	-3	1/6	-2/3	-1	1/6	XQ_L^i
	3	$(\overline{3}, \overline{2}, 1, 1, 1, 1)$	0	-1/3	3	-1/6	2/3	1	-1/6	$\overline{XQ}_L^i$
$a2b'$	3	$(1, 2, 1, 1, 1, 1)$	0	-1	-3	-1/2	-1	0	-1/2	XL_L^i
	3	$(1, \overline{2}, 1, 1, 1, 1)$	0	1	3	1/2	1	0	1/2	$\overline{XL}_L^i$
$a1c1'$	3	$(3, 1, 1, 1, 1, 1)$	1/2	1/3	-3	2/3	-2/3	-1	2/3	XU_R^i
	3	$(\overline{3}, 1, 1, 1, 1, 1)$	-1/2	-1/3	3	-2/3	2/3	1	-2/3	$\overline{XU}_R^i$
$a1c2'$	3	$(3, 1, 1, 1, 1, 1)$	-1/2	1/3	-3	-1/3	-2/3	-1	-1/3	XD_R^i
	3	$(\overline{3}, 1, 1, 1, 1, 1)$	1/2	-1/3	3	1/3	2/3	1	1/3	$\overline{XD}_R^i$
$a2c1'$	3	$(1, 1, 1, 1, 1, 1)$	1/2	-1	-3	0	-1	0	0	XN_R^i
	3	$(1, 1, 1, 1, 1, 1)$	-1/2	1	3	0	1	0	0	$\overline{XN}_R^i$
$a2c2'$	3	$(1, 1, 1, 1, 1, 1)$	-1/2	-1	-3	-1	-1	0	-1	XE_R^i
	3	$(1, 1, 1, 1, 1, 1)$	1/2	1	3	1	1	0	1	$\overline{XE}_R^i$
$bc1$	6	$(1, 2, 1, 1, 1, 1)$	-1/2	0	0	-1/2	0	0	-1/2	H_d^i
	6	$(1, \overline{2}, 1, 1, 1, 1)$	1/2	0	0	1/2	0	0	1/2	$\overline{H}_d^i$
$bc2$	6	$(1, 2, 1, 1, 1, 1)$	1/2	0	0	1/2	0	0	1/2	H_u^i
	6	$(1, \overline{2}, 1, 1, 1, 1)$	-1/2	0	0	-1/2	0	0	-1/2	$\overline{H}_u^i$
$bc1'$	1	$(1, \overline{2}, 1, 1, 1, 1)$	-1/2	0	4	-1/2	1	1	-1/2	$\overline{\mathcal{H}}_1$
	1	$(1, 2, 1, 1, 1, 1)$	1/2	0	-4	1/2	-1	-1	1/2	$\mathcal{H}_1$
$bc2'$	1	$(1, \overline{2}, 1, 1, 1, 1)$	1/2	0	4	1/2	1	1	1/2	$\overline{\mathcal{H}}_2$
	1	$(1, 2, 1, 1, 1, 1)$	-1/2	0	-4	-1/2	-1	-1	-1/2	$\mathcal{H}_2$

where

$$\begin{aligned} U(1)_Y &= \tfrac{1}{6}\left[U(1)_{a1} - 3U(1)_{a2} + 3U(1)_{c1} - 3U(1)_{c2}\right] \\ &= \tfrac{1}{2}U(1)_{B-L} + U(1)_{I3R}, \end{aligned} \tag{7}$$

and

$$\begin{aligned} U(1)_B &= \tfrac{1}{4}[U(1)_{B-L} + U(1)_{3B+L}] \\ &= -[\tfrac{1}{6}U(1)_{a1} + \tfrac{1}{2}(U(1)_{a2} + U(1)_b + U(1)_{c1} + U(1)_{c2} + 3U(1)_d)]. \end{aligned} \tag{8}$$

Alternatively, the gauge symmetry may also be broken by assigning VEVs to the vector-like fields XN_R^i and $\overline{XN}_R^i$ in the a_2c_1 sector shown in Table V. The gauge symmetry in this case is then

$$[SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_L]_{observable} \times \left[U(1)_{I3V} \times USp(2)^2\right]_{hidden} \tag{9}$$

where

$$U(1)_L = \tfrac{1}{4}[-3U(1)_{B-L} + U(1)_{3B+L}]. \tag{10}$$

From Table II, it may be seen that $U(1)_B$ and $U(1)_L$ count baryon number and lepton respectively for the chiral fields, although this is not the case for the vector-like fields. The $U(1)_Y \times U(1)_B \times U(1)_{I3V}$ and $U(1)_Y \times U(1)_L \times U(1)_{I3V}$ gauge symmetries may also be broken by assigning VEVs to some of the vector-like singlet fields $\varphi 21_i$, $\varsigma 21$, $\psi 12$, and $\psi 21$ shown in Table VI. After this breaking, one anomaly-free linear combination remains:

$$U(1)_{Y'} = \frac{1}{6}\left[U(1)_{a1} - 3U(1)_{a2} + 3U(1)_{c1} - 3U(1)_{c2} + 3U(1)_{d1} - 3U(1)_{d2}\right]. \tag{11}$$

The VEVs assigned to the vector-like singlets may be string scale or just below the string scale, so that below the string scale the gauge symmetry is

$$[SU(3)_C \times SU(2)_L \times U(1)_{Y'}]_{observable} \times \left[USp(2)^2\right]_{hidden}, \tag{12}$$

with $U(1)_{Y'}$ being identified with the SM hypercharge. We will further assume that all exotic matter, shown in Table III, may become massive, as shown in Ref. [41]. The resulting low-energy field content is shown in Tables II and along with their charges under $U(1)_{I3R}$, $U(1)_{B-L}$, $U(1)_{3B+L}$, $U(1)_Y$, $U(1)_{I3V}$, $U(1)_B$, and $U(1)_L$. It should be noted that there are several fields present which are SM singlets.

Finally, it is possible to calculate the gauge couplings at the string scale. For this model, it is found that the tree-level gauge couplings are unified at the string scale:

$$g_s^2 = g_w^2 = \frac{5}{3}g_Y^2 = 2g_{Y'}^2, \tag{13}$$

where the unification with $g_{Y'}^2$ is non-canonical [39]. Moreover, the hidden sector gauge groups $USp(2)_3$ and $USp(2)_4$ will become strongly coupled near the string scale, thus decoupling matter charged under these groups [39].

III. GAUGE COUPLING UNIFICATION AND VECTOR-LIKE MATTER

Vector-like matter appears in intersecting/magnetized D-brane models on toroidal orientifolds between stacks of D-branes which do not intersect. The mass of such vector-like

TABLE V. The vector-like quarks and $SU(2)_L$ doublets charged under $U(1)_{I3V}$, their multiplicities and quantum numbers under the gauge symmetry $[SU(3)_C \times SU(2)_L \times U(1)_Y]_{observable}\left[U(1)_{I3V} \times USp(2)^2\right]_{hidden}$ and their charges under different $U(1)$ groups. Here a, b, c, etc. refer to different stacks of D-branes.

	Mult.	Quantum Number	Q_{I3R}	Q_{B-L}	Q_{3B+L}	Q_Y	$Q_{I_{3V}}$	Q_B	Q_L	$Q_{Y'}$	Field
$a1d1$	2	$(3,1,1,1,1,1)$	0	1/3	5	1/6	1/2	4/3	1	2/3	$\varphi 11_i$
	2	$(\overline{3},1,1,1,1,1)$	0	-1/3	-5	-1/6	-1/2	-4/3	-1	-2/3	$\overline{\varphi 11}_i$
$a1d2$	2	$(3,1,1,1,1,1)$	0	1/3	5	1/6	-1/2	4/3	1	-1/3	$\varphi 12_i$
	2	$(\overline{3},1,1,1,1,1)$	0	-1/3	-5	-1/6	1/2	-4/3	-1	1/3	$\overline{\varphi 12}_i$
$a1d1'$	1	$(3,1,1,1,1,1)$	0	1/3	-7	1/6	1/2	-5/3	-2	2/3	$\varsigma 11$
	1	$(\overline{3},1,1,1,1,1)$	0	-1/3	7	-1/6	-1/2	5/3	2	-2/3	$\overline{\varsigma 11}$
$a1d2'$	1	$(3,1,1,1,1,1)$	0	1/3	-7	1/6	-1/2	-5/3	-2	-1/3	$\varsigma 12$
	1	$(\overline{3},1,1,1,1,1)$	0	-1/3	7	-1/6	1/2	5/3	2	1/3	$\overline{\varsigma 12}$
$bd1'$	1	$(1,2,1,1,1,1)$	0	0	-8	0	1/2	-2	-2	1/2	ξ_1
	1	$(1,\overline{2},1,1,1,1)$	0	0	8	0	-1/2	2	2	-1/2	$\overline{\xi}_1$
$bd2'$	1	$(1,2,1,1,1,1)$	0	0	-8	0	-1/2	-2	-2	-1/2	ξ_2
	1	$(1,\overline{2},1,1,1,1)$	0	0	8	0	1/2	2	2	1/2	$\overline{\xi}_2$

states depends upon the separation between the stacks of D-branes in the internal space. As such, it is generically massive. Only stacks of D-branes which are directly on top of one another have massless vector-like states between them. In the model studied in Section II, the toroidal orientifold consist of a six-torus which is factorizable, $\mathbf{T}^6 = \mathbf{T}^2 \times \mathbf{T}^2 \times \mathbf{T}^2$. If two stacks of D-branes are parallel on one two-torus, then vector-like matter appears in the bifundamental representation of the gauge groups within the world-volume of each stack. If the two stacks are not separated on the two torus on which they are parallel, the vector-like multiplets are massless. However, these states become massive if the stacks are separated.

The most straightforward way to obtain the diphoton excesses is with a SM singlet scalar with a mass $\sim$ 750 GeV coupled to vector-like quarks. The coupling to vector-like quarks is necessary in order to produce the scalar via gluon fusion and to allow the decay of the scalar into diphotons. The requisite vector-like quarks are indeed present in the model. In particular, the vector-like quarks and leptons in the ab' and ac' sector fill $\mathbf{16}$ and $\overline{\mathbf{16}}$ spinorial

TABLE VI. The vector-like singlets charged under $U(1)_{I3V}$,their multiplicities and quantum numbers under the gauge symmetry $[SU(3)_C \times SU(2)_L \times U(1)_Y]_{observable} \left[U(1)_{I3V} \times USp(2)^2\right]_{hidden}$ and their charges under different $U(1)$ groups. Here a, b, c, etc. refer to different stacks of D-branes.

	Mult.	Quantum Number	Q_{I3R}	Q_{B-L}	Q_{3B+L}	Q_Y	$Q_{I_{3V}}$	Q_B	Q_L	$Q_{Y'}$	Field
$a2d1$	2	$(1,1,1,1,1,1)$	0	-1	5	-1/2	1/2	1	1/2	0	$\varphi21_i$
	2	$(1,1,1,1,1,1)$	0	1	-5	1/2	-1/2	-1	-1/2	0	$\overline{\varphi21}_i$
$a2d2$	2	$(1,1,1,1,1,1)$	0	-1	5	-1/2	-1/2	1	1/2	-1	$\varphi22_i$
	2	$(1,1,1,1,1,1)$	0	1	-5	1/2	1/2	-1	-1/2	1	$\overline{\varphi22}_i$
$a2d1'$	1	$(1,1,1,1,1,1)$	0	-1	-7	-1/2	1/2	-2	-1	0	$\varsigma21$
	1	$(1,1,1,1,1,1)$	0	1	7	1/2	-1/2	2	1	0	$\overline{\varsigma21}$
$a2d2'$	1	$(1,1,1,1,1,1)$	0	-1	-7	-1/2	-1/2	-2	-1	-1	$\varsigma22$
	1	$(1,1,1,1,1,1)$	0	1	7	1/2	1/2	2	1	1	$\overline{\varsigma22}$
$c1d1'$	1	$(1,1,1,1,1,1)$	1/2	0	-8	1/2	1/2	-2	-2	1	$\psi11$
	1	$(1,1,1,1,1,1)$	-1/2	0	8	-1/2	-1/2	2	2	-1	$\overline{\psi11}$
$c1d2'$	1	$(1,1,1,1,1,1)$	1/2	0	-8	1/2	-1/2	-2	-2	0	$\psi12$
	1	$(1,1,1,1,1,1)$	-1/2	0	8	-1/2	1/2	2	2	0	$\overline{\psi12}$
$c2d1'$	1	$(1,1,1,1,1,1)$	-1/2	0	-8	-1/2	1/2	-2	-2	0	$\psi21$
	1	$(1,1,1,1,1,1)$	1/2	0	8	-1/2	-1/2	2	2	0	$\overline{\psi21}$
$c2d2'$	1	$(1,1,1,1,1,1)$	-1/2	0	-8	-1/2	-1/2	-2	-2	-1	$\psi22$
	1	$(1,1,1,1,1,1)$	1/2	0	8	1/2	1/2	2	2	1	$\overline{\psi22}$

representations of SO(10), or equivalently $\mathbf{5}+\overline{\mathbf{5}}+\mathbf{10}+\overline{\mathbf{10}}+\mathbf{1}+\overline{\mathbf{1}}$ of SU(5). For example,

$$\begin{aligned} \mathbf{5}^i+\overline{\mathbf{5}}^i &= \left\{(XD^i_R,\overline{XD}^i_R),(XL^i_L,\overline{XL}^i_L)\right\}, \\ \mathbf{10}^i+\overline{\mathbf{10}}^i &= \left\{(XQ^i_L,\overline{XQ}^i_L),(XU^i_R,\overline{XU}^i_R),(XE^i_R,\overline{XE}^i_R)\right\}, \\ \mathbf{1}^i+\overline{\mathbf{1}}^i &= \left\{XN^i_R,\overline{XN}^i_R\right\}. \end{aligned} \tag{14}$$

It is well-known that gauge coupling unification may be preserved at the 1-loop level if the extra matter comes in complete representations of $SU(5)$. However, at the 2 or 3-loop level, a Landau pole may appear. This restricts the number of $SU(5)$ multiplets which may remain light to either one $(\mathbf{10}+\overline{\mathbf{10}})$ or three copies of $(\mathbf{5}+\overline{\mathbf{5}})$. In addition, any number of SM singlets may be present in the light spectrum. Generically, there are many more vector-like states in the spectrum, as can be seen from Tables IV,V, and VI. Thus, in order to preserve gauge coupling unification, many of these states must obtain string-scale masses, while simultaneously maintaining light masses for either of the two cases stated above. Recalling that vector-like matter appears between stacks of D-branes which are parallel on one two-torus, the masses of these vector-like states depend on the separation between the stacks on the two-torus on which they are parallel. So, it is possible to give string-scale masses to some of the vector-like states by separating the two stacks of D-branes on the two-torus on which they are parallel and where these vector-like states are localized.

Is it possible to choose the position of the stacks of D-branes so that only the $\mathbf{5}+\overline{\mathbf{5}}$ or $\mathbf{10}+\overline{\mathbf{10}}$ fields remain light? The answer to this question is yes. As an example, let us consider just the fields in the ab' and ac' sectors, where it should be noted that these stacks of D-branes are parallel on the third two-torus. Thus, for example the fields $(XQ^i_L, \overline{XQ}^i_L)$ may become massive by separating stacks $a1$ and b' on the third torus. Similarly, the fields $(XU^i_R, \overline{XU}^i_R)$ may become massive by separating stacks $a1$ and $c1'$, while $(XE^i_R, \overline{XE}^i_R)$ become massive if stacks $a2$ and $c2'$ are separated on the third two-torus. The fields $(XD^i_R, \overline{XD}^i_R)$ may remain massless if stacks $a1$ and $c2'$ overlap on the third two-torus, and the same is true for $(XL^i_L, \overline{XL}^i_L)$ if stacks $a2$ and b' overlap on the third two-torus. Thus with this configuration, only the fields in the $\mathbf{5}^i+\overline{\mathbf{5}}^i$ and $\mathbf{1}^i+\overline{\mathbf{1}}^i$ (with $i=1\ldots3$) representations of Eq. 14 remain light.

In addition, there are additional vector-like quarks and leptons in the spectrum in the ad and ad' sectors, as shown in Tables V and VI. These fields may also become massive by separating the relevant stacks of D-branes on the two-torus where they are parallel. For example, stacks $a1$ and $a2$ are parallel with stacks $d1$ and $d2$ on the third two-torus, and so the vector-like matter in these sectors may be eliminated if these stacks are separated on the third two-torus. Similarly, stacks $a1$ and $a2$ are parallel with stacks $d1'$ and $d2'$ on the second two-torus, and the vector-like matter in this sector may be eliminated by displacing these stacks on the second two-torus.

On the other hand, it must also be kept in mind that splitting the stacks in this fashion may affect the Higgs sectors, which are also vector-like. For example the Higgs fields H_u and H_d arise from vector-like matter in the $bc1$ and $bc2$ sectors respectively. Stackss $a1$ and $c2'$ must overlap to keep $(XE^i_R, \overline{XE}^i_R)$ light while stacks $a1$ and b' must be separated to give masses to the fields $(XQ^i_L, \overline{XQ}^i_L)$ such that the light vector-like fields may be placed in $\mathbf{5}+\overline{\mathbf{5}}$ representations. However, this implies that stacks b and $c2$ must be separated and thus that the Higgs field H_d is not present in the light spectrum. Clealy, then it is not possible to eliminate the fields in the $\mathbf{10}+\overline{10}$ while keeping the fields in the $\mathbf{5}+\overline{5}$ without also eliminating some of the Higgs field. Note that this constraint only applies to those fields in the ab' and ac' sectors. In the next section, we shall also consider the vector-like matter in the ad and ad' sectors. We shall find that it is possible to have just vector-like matter in the $\mathbf{5}+\overline{5}$ representation without eliminating some of the Higgs fields as a by-product.

IV. MSSM SINGLETS COUPLED TO VECTOR-LIKE MATTER

We have seen in the previous section that the model contains vector-like quarks and leptons in the ab' and ac' sectors which may be grouped into complete representations of $SU(5)$. Furthermore, by displacing the stacks of D-branes, it is possible to eliminate some of these fields from the spectrum so that only three copies of $\mathbf{5}+\overline{\mathbf{5}}$ remain light. In this way, gauge coupling unification may be maintained while also avoiding the Landau pole problem. However, upon inspection, none of these vector-like quarks and leptons appear to have Yukawa couplings with any of the singlet fields present in the model. Therefore, they may not be involved in producing the diphoton excesses. However, this is not the case if we examine the vector-like quarks and leptons present in the ad and ad' sectors.

Let us turn our attention to the SM singlet fields X_{c1d2} and X_{c2d1} shown in Table III. These fields have Yukawa couplings with the vector-like quarks in the ad and ad' sectors:

$$W_3 = c_1^i \cdot \overline{\psi 12} \cdot X_{c1d2} \cdot \varphi 12^i \cdot \varsigma 12 + c_2^i \cdot \overline{\psi 21} \cdot X_{c2d1} \cdot \varphi 11^i \cdot \varsigma 11, \tag{15}$$

where $\langle\overline{\psi 12}\rangle$ and $\langle\overline{\psi 21}\rangle$ are near the string scale M_{St} as discussed in Section I. In addition, there are Yuakwa couplings between these singlets and some of the vector-like doublets:

$$W_2 = k_1 X_{c1d2}\overline{\mathcal{H}}_1\xi_2 + k_2 X_{c2d1}\overline{H}_2\xi_1. \tag{16}$$

Finally, there are additional Yukawa couplings between X_{c1d2} and X_{c2d1} and the singlet fields in the $a2d$ and $a2d'$ sectors:

$$W_1 = d_1^i \cdot \overline{\psi 12} \cdot X_{c1d2} \cdot \varphi 22^i \cdot \varsigma 22 + d_2^i \cdot \overline{\psi 21} \cdot X_{c2d1} \cdot \varphi 21^i \cdot \varsigma 21. \tag{17}$$

Thus, the required couplings of the SM singlet fields to vector-like quarks are present so that the singlet fields may be produced by gluon fusion and decay to diphotons. In addition, the singlet fields have additional couplings to other doublet and singlet fields. Note that the vector-like quarks involved in these couplings combined with additional vector-like matter may be placed in complete representations of $SU(5)$ by replacing the right-handed vector-like quarks, leptons, and singlets in Eq. 14 with some of the quarks, leptons and singlets of

Tables V and VI. For example, making the interchanges

$$\begin{aligned} (XD_R^i, \overline{XD}_R^i) &\to (\varphi 12^i, \overline{\varphi 12}^i), \\ (XU_R^i, \overline{XU}_R^i) &\to (\varphi 11^i, \overline{\varphi 11}^i), \\ (XE_R^i, \overline{XE}_R^i) &\to (\varphi 22^i, \overline{\varphi 22}^i), \\ (XN_R^i, \overline{XN}_R^i) &\to (\varphi 21^i, \overline{\varphi 21}^i), \\ (XL_L^i, \overline{XL}_L^i) &\to (\mathcal{H}^i, \overline{\mathcal{H}}^i), \end{aligned} \tag{18}$$

we have

$$\begin{aligned} \mathbf{5}^i + \overline{5}^i &= \left\{ (\varphi 12_i, \overline{\varphi 12}^i), (\mathcal{H}^i, \overline{\mathcal{H}}^i) \right\}, \\ \mathbf{10}^i + \overline{10}^i &= \left\{ (XQ_L^i, \overline{XQ}_L^i), (\varphi 11_i^i, \overline{\varphi 11}^i), (\varphi 22^i, \overline{\varphi 22}^i) \right\}, \\ \mathbf{1}^i + \overline{\mathbf{1}}^i &= \left\{ \varphi 21^i, \overline{\varphi 21}^i \right\}. \end{aligned} \tag{19}$$

where $i = 1 \ldots 2$, and making the interchanges

$$\begin{aligned} (XD_R^3, \overline{XD}_R^3) &\to (\varsigma 12, \overline{\varsigma 12}), \\ (XU_R^3, \overline{XU}_R^3) &\to (\varsigma 11, \overline{\varsigma 11}), \\ (XE_R^3, \overline{XE}_R^3) &\to (\varsigma 22, \overline{\varsigma 22}), \\ (XN_R^3, \overline{XN}_R^3) &\to (\varsigma 21, \overline{\varsigma 21}), \\ (XL_L^3, \overline{XL}_R^3) &\to (\xi_1, \overline{\xi_1}), \end{aligned} \tag{20}$$

we have

$$\begin{aligned} \mathbf{5} + \overline{\mathbf{5}} &= \left\{ (\varsigma 12, \overline{\varsigma 12}), (\xi_1, \overline{\xi_1}) \right\}, \\ \mathbf{10} + \overline{\mathbf{10}} &= \left\{ (XQ_L^3, \overline{XQ}_L^3), (\varsigma 11, \overline{\varsigma 11}), (\varsigma 22, \overline{\varsigma 22}) \right\}, \\ \mathbf{1} + \overline{\mathbf{1}} &= \left\{ \varsigma 21, \overline{\varsigma 21} \right\}. \end{aligned} \tag{21}$$

In order to have just three $\mathbf{5}+\overline{\mathbf{5}}$ multiplets + additional singlets in the light spectrum, the fields grouped into the $\mathbf{10}+\overline{\mathbf{10}}$ multiplets must become massive, as well as any additional vector-like states. To eliminate the vectorlike fields in the ab' and ac' sectors from the light spectrum, we must require that stacks $a1$ and $a2$ are separated from stacks b', $c1'$, and $c2'$ on the third two-torus. Futhermore, to eliminate the fields in the above $\mathbf{10}+\overline{\mathbf{10}}$, we require that stack $a1$ be separated from stack $d1$ on the third two-torus and from $d1'$ on the second two-torus. Let us also require that stack $a2$ be separated from both stacks $d1$, $d1'$, $d2$ and $d2'$.

To keep the fields in the above $\mathbf{5}+\overline{\mathbf{5}}$ representations light, we require that stack $a1$ overlap with stack $d2$ on the third two-torus and with stack $d2'$ on the second two torus. In addition, we require that stack b overlap stacks $c1$, $c1'$, $c2$, and $c2'$ as well as stacks $d1'$. This configuration assures that the Higgs fields H_u^i and H_d^i are present in the spectrum. In addition, stack b may not overlap stack $d1'$ since stack $a1$ overlaps stack $d1$. Since stack $a1$ and b' are separated, this implies that stacks b and $d1'$ are also separated.

Then, the fields in the light spectrum consist of the following fields with quantum numbers under the $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge symmetry shown:

$$\begin{aligned}(XD_{1,2}, XD^c_{1,2}) &\equiv 2\times(\varphi12, \overline{\varphi12}) = 2\times\left\{(\overline{\mathbf{3}},\mathbf{1},\frac{\mathbf{1}}{\mathbf{3}}) + (\mathbf{3},\mathbf{1},-\frac{\mathbf{1}}{\mathbf{3}})\right\}, \\ (XD_3, XD^c_3) &\equiv 1\times(\varsigma12, \overline{\varsigma12}) = \left\{(\overline{\mathbf{3}},\mathbf{1},\frac{\mathbf{1}}{\mathbf{3}}) + (\mathbf{3},\mathbf{1},-\frac{\mathbf{1}}{\mathbf{3}})\right\}, \\ (XL_{L_{1,2}}, XL^c_{L_{1,2}}) &\equiv 2\times(\mathcal{H}_1, \overline{\mathcal{H}_1}) = \left\{(\mathbf{1},\overline{\mathbf{2}},-\frac{\mathbf{1}}{\mathbf{2}}) + (\mathbf{1},\mathbf{2},\frac{\mathbf{1}}{\mathbf{2}})\right\}, \\ (XL_{L_3}, XL^c_{L_3}) &\equiv 1\times(\xi_1, \overline{\xi_1}) = \left\{(\mathbf{1},\overline{\mathbf{2}},-\frac{\mathbf{1}}{\mathbf{2}}) + (\mathbf{1},\mathbf{2},\frac{\mathbf{1}}{\mathbf{2}})\right\},\end{aligned} \tag{22}$$

plus additional singlets. As shown in Eq. 15, the vector-like quarks $(\varphi12^{1,2}, \overline{\varphi12}^{1,2})$ and $(\varsigma12, \overline{\varsigma12})$ have Yukawa couplings with the singlet field $S \equiv X_{c1d2}$. Thus, this singlet field may be produced via loops involving these vector-like quarks as well as decay to diphotons via gluon fusion. In addition, the singlet field S has couplings to the doublets (XL_L, XL_L^c) to which it may also decay.

Using the notation of Eq. 22, the superpotential for the extra vector-like states and the singlet S is

$$\begin{aligned} W = {} & \lambda_D S XD XD^c + \lambda_L S XL XL^c \\ & + M_{XD} XD XD^c + M_{XL} XL XL^c \ . \end{aligned} \tag{23}$$

The corresponding supersymmetry breaking soft terms are

$$\begin{aligned} V_{\rm soft} = {} & \widetilde{M}^2_{XD}(|\widetilde{XD}|^2 + |\widetilde{XD}^c|^2) + \widetilde{M}^2_{XL}(|\widetilde{XL}|^2 + |\widetilde{XL}^c|^2) \\ & - \Big(\lambda_D A_D S \widetilde{XD}\widetilde{XD}^c \\ & + \lambda_L A_L S \widetilde{XL}\widetilde{XL}^c \\ & + B_{XD} M_{XD} XD XD^c + B_{XL} M_{XL} XL XL^c + \text{H.C.}\Big) \ . \end{aligned} \tag{24}$$

V. THE DIPHOTON EXCESSES

The 750 GeV diphoton production cross-sections observed by the CMS collaboration are $\sigma(pp \to S \to \gamma\gamma) = 0.5 \pm 0.6$ fb at $\sqrt{s} = 8$ TeV [47]and 6 ± 3 fb at $\sqrt{s} = 13$ TeV [2], while

TABLE VII. Decay widths and production cross-sections for a total decay width of $\Gamma = 5$ GeV for some sample points. All masses and decay widths are in GeV. The cross-sections are in femtobarns (fb). The $\mathrm{Br}_{\mathrm{DM}}$ represents the branching ratio allocated to dark matter. For simplicity, we assume here that $M_{XE} = M_{XN}$ and $\lambda_E = \lambda_N$.

$\Gamma = 5$ GeV																		
M_{XD}	M_{XE}	$M_{\widetilde{XD}}$	$M_{\widetilde{XE}}$	$\widetilde{M}_{XD}$	$\widetilde{M}_{XE}$	λ_D	λ_E	A_D	A_E	$\Gamma_{\gamma\gamma}$	Γ_{gg}	Γ_{XE+XN}	$\mathrm{Br}_{\gamma\gamma}$	Br_{gg}	Br_{XE+XN}	$\mathrm{Br}_{\mathrm{DM}}$	$\sigma_{\gamma\gamma}^{8\ \mathrm{TeV}}$	$\sigma_{\gamma\gamma}^{13\ \mathrm{TeV}}$
1200	255	2050	400	1662	308	0.65	0.351	3600	1530	0.0037	0.185	1.45	0.0007	0.037	0.290	0.672	0.38	1.77
1350	225	1750	330	1114	241	0.55	0.351	4050	1350	0.0077	0.138	1.88	0.0015	0.028	0.376	0.594	0.60	2.77
1200	265	1800	310	1342	161	0.60	0.351	3600	1590	0.0083	0.178	1.30	0.0017	0.036	0.260	0.702	0.82	3.83
1400	235	1750	330	1050	232	0.70	0.351	4200	1410	0.0081	0.218	1.74	0.0016	0.044	0.348	0.607	0.99	4.58
1050	225	1800	300	1462	198	0.70	0.351	3150	1350	0.0072	0.284	1.88	0.0014	0.057	0.376	0.565	1.14	5.32
1000	215	1450	330	1050	250	0.65	0.351	3000	1290	0.0075	0.319	2.02	0.0015	0.064	0.404	0.530	1.34	6.22
1000	255	1500	330	1118	209	0.65	0.351	3000	1530	0.0088	0.307	1.45	0.0018	0.061	0.290	0.647	1.50	7.00

TABLE VIII. Decay widths and production cross-sections for a total decay width of $\Gamma = 45$ GeV for some sample points. All masses and decay widths are in GeV. The cross-sections are in femtobarns (fb). The $\mathrm{Br}_{\mathrm{DM}}$ represents the branching ratio allocated to dark matter. For simplicity, we assume here that $M_{XE} = M_{XN}$ and $\lambda_E = \lambda_N$.

$\Gamma = 45$ GeV																		
M_{XD}	M_{XE}	$M_{\widetilde{XD}}$	$M_{\widetilde{XE}}$	$\widetilde{M}_{XD}$	$\widetilde{M}_{XE}$	λ_D	λ_E	A_D	A_E	$\Gamma_{\gamma\gamma}$	Γ_{gg}	Γ_{XE+XN}	$\mathrm{Br}_{\gamma\gamma}$	Br_{gg}	Br_{XE+XN}	$\mathrm{Br}_{\mathrm{DM}}$	$\sigma_{\gamma\gamma}^{8\ \mathrm{TeV}}$	$\sigma_{\gamma\gamma}^{13\ \mathrm{TeV}}$
1200	255	1750	340	1274	225	0.70	0.351	3600	1530	0.0090	0.250	1.45	0.00020	0.0055	0.032	0.962	0.14	0.65
950	255	1300	300	887	158	0.65	0.351	2850	1530	0.0079	0.380	1.45	0.00018	0.0085	0.032	0.959	0.19	0.87
1100	255	1450	350	945	240	0.70	0.351	3300	1530	0.0092	0.337	1.45	0.00021	0.0075	0.032	0.960	0.19	0.90
1000	265	1350	330	907	197	0.70	0.351	3000	1590	0.0091	0.401	1.30	0.00020	0.0089	0.029	0.962	0.23	1.05
900	265	1250	340	867	213	0.65	0.351	2700	1590	0.0094	0.420	1.30	0.00021	0.0093	0.029	0.962	0.25	1.14
950	265	1300	350	887	229	0.70	0.351	2850	1590	0.0097	0.441	1.30	0.00021	0.0098	0.029	0.961	0.26	1.23
800	265	1200	360	894	244	0.70	0.351	2400	1590	0.0099	0.580	1.30	0.00022	0.0129	0.029	0.958	0.36	1.67

the ATLAS collaboration observed $\sigma(pp \to S \to \gamma\gamma) = 0.4 \pm 0.8$ fb at $\sqrt{s} = 8$ TeV [48] and 10 ± 3 fb at $\sqrt{s} = 13$ TeV [1]. Replicating the strategy of Ref. [15], we constrain the total decay width to $\Gamma \sim 5 - 45$ GeV. To reproduce the observed production cross-sections, we constrain the model using $\Gamma_{\gamma\gamma}\Gamma_{gg}/M_S^2 \gtrsim 10^{-9}$.

The effective loop-level couplings amongst the Standard Model gauge bosons and scalar S are given by

$$-\mathcal{L} = \frac{S}{M_S}\left[\kappa_{EM} F_{\mu\nu}^{EM} F^{EM\mu\nu} + \kappa_3 G_{\mu\nu}^{a} G^{\mu\nu\, a}\right] \tag{25}$$

where $F_{\mu\nu}^{EM}$ and $G_{\mu\nu}^{a}$ are the photon and gluon field strength tensors, respectively, with $a = 1,\ 2, ..8$. The effective operators are represented by κ_{EM} and κ_3, which are written as

$$\kappa_{EM} = \frac{\alpha_{EM}}{4\pi}\left[\sum_f \frac{\lambda_f M_S}{M_f} Q_f^2 N_{EM}^f F_f + \sum_{\widetilde{f}} \frac{\lambda_f A_f M_S}{M_{\widetilde{f}}^2} Q_{\widetilde{f}}^2 N_{EM}^f F_{\widetilde{f}}\right] \tag{26}$$

$$\kappa_3 = \frac{\alpha_3}{4\pi}\left[\sum_f \frac{\lambda_f M_S}{M_f} N_3^f F_f + \sum_{\widetilde{f}} \frac{\lambda_f A_f M_S}{M_{\widetilde{f}}^2} N_3^f F_{\widetilde{f}}\right] \tag{27}$$

where $N_{EM}^{XD} = 3$, $N_{EM}^{XE} = 1$, and $N_3^{XD} = 1$, and the functions F_f and $F_{\widetilde{f}}$ are expressed as

$$F_f = 2\chi\left[1 + (1-\chi)f(\chi)\right] \tag{28}$$

$$F_{\widetilde{f}} = \chi\left[-1 + \chi f(\chi)\right] \tag{29}$$

with the function χ denoted by

$$\chi = 4\frac{M_{f/\widetilde{f}}^2}{M_S^2} \tag{30}$$

The triangle loop functions $f(\chi)$ are defined here as

$$f(\chi) = \begin{cases} \arcsin^2[\sqrt{\chi^{-1}}] & \text{if } \chi \geq 1 \\ -\frac{1}{4}\left[\ln\frac{1+\sqrt{1-\chi}}{1-\sqrt{1-\chi}} - i\pi\right]^2 & \text{if } \chi < 1. \end{cases} \tag{31}$$

The diphoton and digluon decay widths in $\mathcal{F}$-$SU(5)$ are computed from

$$\Gamma_{\gamma\gamma} = \frac{|\kappa_{EM}|^2}{4\pi}M_S \tag{32}$$

$$\Gamma_{gg} = \frac{2\,|\kappa_3|^2}{\pi}M_S \tag{33}$$

The diphoton production cross-section is calculated from

$$\sigma(pp \to S \to \gamma\gamma) = K\frac{C_{gg}\Gamma(S \to gg)\Gamma(S \to \gamma\gamma)}{s\Gamma M_S} \tag{34}$$

where K is the QCD K-factor, Γ is the total decay width, $\sqrt{s}$ is the proton-proton center of mass energy, and C_{gg} is the dimensionless partonic integral computed for an $M_S = 750$ GeV resonance, yielding $C_{gg} = 174$ at $\sqrt{s} = 8$ TeV and $C_{gg} = 2137$ at $\sqrt{s} = 13$ TeV [49]. We use the gluon fusion K-factor of 1.98.

We construct our intersecting D-brane model with the (XD,XD^c) and (XL,XL^c) vector-like particles, implementing three copies of the $(\mathbf{5},\overline{\mathbf{5}})$. For the calculations, we decompose the (XL, XL^c) multiplet into its (XE, XE^c) and (XN, XN^c) components. Given the null XN electric charge $Q_{XN} = 0$, no constraints can be placed on M_{XN}, λ_N, or A_N in the model via the production cross-section calculations, so for simplicity we set $M_{XN} = M_{XE}$ and $\lambda_N = \lambda_E$ when computing the decay of the scalar S directly to the XN multiplet. The multiplets (XD, XD^c) and (XE, XE^c) participate in the $S \to \gamma\gamma$ loop diagrams and only (XD, XD^c) in the $S \to gg$ loops, hence there are 8 free parameters in the effective operators κ_{EM} and κ_3 consisting of the Yukawa couplings λ_f, trilinear A term couplings A_f, fermionic component masses M_f, and scalar component masses $M_{\widetilde{f}}$. In total, there are 10 parameters to compute:

$$M_{XD},\ M_{XE},\ M_{\widetilde{XD}},\ M_{\widetilde{XE}},\ \widetilde{M}_{XD},\ \widetilde{M}_{XE},\ \lambda_D,\ \lambda_E,\ A_D,\ A_E \tag{35}$$

though the supersymmetry breaking soft terms $\widetilde{M}_{XD}$ and $\widetilde{M}_{XE}$ can be trivially computed from the fermionic and scalar components using the following relations

$$M^2_{\widetilde{XD}} = M^2_{XD} + \widetilde{M}^2_{XD} \quad (36)$$

$$M^2_{\widetilde{XE}} = M^2_{XE} + \widetilde{M}^2_{XE} \quad (37)$$

The freedom on the 8 free-parameters engenders a large D-brane model parameter space. We treat the fermionic component of the vector-like particle and its soft supersymmetry breaking term independently, such that $M_f \neq \widetilde{M}_f$. Recent constraints at the LHC on vector-like B quarks [50] provide lower limits of around 735 GeV for the XD multiplet, allowing a light XE multiplet, which contributes to invisible branching fractions when $M_{XE} < 375$ GeV. For the large Yukawa couplings λ_D and λ_E at the unification scale, there exist the quasi focus points with $\lambda_D \simeq 0.70$ and $\lambda_E \simeq 0.351$ at low energy, respectively [14]. Constraints are placed on the A terms of $A_D \lesssim 3M_D$ and $A_E \lesssim 6M_E$ to preclude premature breaking of the $SU(3)_C \times U(1)_{EM}$ gauge symmetry. We provide some sample benchmark points in TABLE VII and TABLE VIII to illustrate that the model can accommodate the observed diphoton cross-section. Note that the cross-sections $\sigma^{8\text{ TeV}}_{\gamma\gamma}$ and $\sigma^{13\text{ TeV}}_{\gamma\gamma}$ in TABLES VII - VIII display a gain of 4.65 from 8 TeV to 13 TeV. We use values of the coupling constants at the M_Z scale in our calculations of $\alpha_3 = 0.1185$ and $\alpha_{EM} = 128.91^{-1}$.

The diphoton and digluon decay modes are only a small portion of the total decay width, as shown in TABLES VII - VIII. These practical constraints on the diphoton and digluon decay widths allocate the majority of the total width to the other decay channels, for instance, to dark matter invisibly and/or $S \to XEXE^c$ and $S \to XNXN$, kinematically allowed for both M_{XE} and M_{XN} less than 375 GeV, which is reflected in the Γ_{XE+XN} decay width and Br_{XE+XN} branching ratio in TABLES VII - VIII. Specifically, the decay width for the decay into pairs of fermions $S \to f\overline{f}$ is given by

$$\Gamma(S \to f\overline{f}) = \frac{1}{16\pi} M_S \lambda_f^2 \left(1 - \frac{4M_f^2}{M_S^2}\right)^{3/2} . \quad (38)$$

where we take $M_{XE} = M_{XN}$ and $\lambda_E = \lambda_N$ in the calculations.

VI. CONCLUSION

We have studied the diphoton excesses near 750 GeV recently reported by the ATLAS and CMS collaborations within the context of a phenomenologically interesting intersecting/magnetized D-brane model on a toroidal orientifold. We have shown that the model contains a SM singlet scalar as well as vector-like quarks and leptons. In addition, we have shown that the singlet scalar has Yukawa couplings with vector-like quarks and leptons such that it may be produced in proton-proton collisions via gluon fusion as well as decay to diphotons through loops involving the vector-like quarks. Moreover, the required vector-like quarks and leptons may appear in complete $SU(5)$ multiplets so that gauge coupling unification may be maintained. In particuar, we showed that we may have three copies of $\mathbf{5} + \overline{\mathbf{5}}$ representations of $SU(5)$ in the light spectrum which are present in the model. Finally, we showed that the diphoton excesses observed by the ATLAS and CMS collaborations may be accommodated.

It should be emphasized that we have obtained these results within the context of a complete, globally consistent string model. This particular model has many interesting phenomenological features such as automatic gauge coupling unification, realistic Yukawa mass matrices for quarks and leptons, and mimimal exotic matter. It is of note that the singlet fields required to explain the diphoton signal arise from this extra matter. In addition, the required vector-like quarks are naturally present in the spectrum. Finally, the Yukawa couplings between the singlet fields and the vector-like quarks are allowed by the global symmetries arising from $U(1)$ factors whose gauge bosons become heavy at the string scale via the Green-Schwarz mechanism, a result which is completely non-trivial.

An interesting question is whether or not is it possible to obtain supersymmetry partner spectra from the model which take into account the light vector-like matter. As mentioned earlier, the vector-like quarks and leptons have Yukawa couplings with the Higgs fields in the model, and thus may raise the Higgs mass by up to a few GeV. This may alleviate the problem with electroweak fine-tuning. In addition, the vector-like matter affects the RGE running of the soft masses. Finally, the soft supersymmetry breaking masses may be calculated at the string scale in the model. Thus, it would be very interesting to study the possible supersymmetry partner spectra obtainable in the model including the extra vector-like matter. We plan to study this in future work.

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