The 750 GeV di-photon LHC excess and extra Z’s in heterotic-string derived models

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Abstract The ATLAS and CMS collaborations recently recorded possible di-photon excess at 750 GeV and a less significant di-boson excess around 1.9 TeV. Such excesses may be produced in heterotic string derived \textit{Z’} models, where the di-photon excess may be connected with the Standard Model singlet scalar responsible for the \textit{Z’} symmetry breaking, whereas the di-boson excess arises from production of the extra vector boson. Additional vector-like states in the string \textit{Z’} model are instrumental to explain the relatively large width of the di-photon events and mandated by anomaly cancellation to be in the vicinity of the \textit{Z’} breaking scale. Wilson line breaking of the non-Abelian gauge symmetries in the string models naturally gives rise to dark matter candidates. Future collider experiments will discriminate between the high-scale heterotic-string models, which preserve the perturbative unification paradigm indicated by the Standard Model data, versus the low scale string models. We also discuss the possibility for the production of the diphoton events with high scale \textit{U(1)}\textsubscript{\textit{Z’}} breaking.

1 Introduction

The Standard Model matter spectrum strongly favours its embedding in SO(10) chiral representations. This unification scenario is further supported by the perturbative logarithmic evolution of the Standard Model parameters; by proton longevity; and by the suppression of left-handed neutrino masses. This picture is reproduced in heterotic-string models \cite{1,2}. The free fermionic formulation \cite{3–5} in particular has given rise to phenomenological three generation models that have been used to explore the unification of gravity and the gauge interactions. These models correspond to \textit{Z}_2 \times \textit{Z}_2 orbifold compactification at special points in the Narain moduli space and utilise discrete Wilson lines to break the non-Abelian gauge symmetry to an SO(10) subgroup \cite{6–9}. The viable models constructed to date include the flipped SU(5) (FSU5) \cite{10}; the standard-like models (SLM) \cite{11–15}; the Pati–Salam models (PS) \cite{16–18} and the left–right symmetric models (LRS) \cite{19,20}, whereas the SU(4) \times SU(2) \times U(1) subgroup did not produce viable model data, versus the low scale string models. We also discuss the possibility for the production of the diphoton events with high scale \textit{U(1)}\textsubscript{\textit{Z’}} breaking.
enhanced to $E_6$. This is the case in the free fermionic model if the different 16 and 10 + 1 representations, which would form a complete 27 representation of $E_6$, are obtained from different fixed points of the underlying $Z_2 \times Z_2$ orbifold [24]. Adding the basis vector that breaks the SO(10) symmetry to the PS subgroup results in split multiplets, but the chiral spectrum still forms complete $E_6$ multiplets, hence rendering $U(1)_C$ anomaly free. We remark that while complete $E_6$ multiplets are sufficient for $U(1)_C$ to be anomaly free, it might not be necessary and alternative possibilities may exist.

The ATLAS and CMS collaborations [31,32] reported recently an excess in the di-photon searches at di-photon invariant mass of 750 GeV. This excess can be attributed to a neutral scalar resonance of 750 GeV mass. Plausible candidates include the SO(10) neutral singlet in the 27 of $E_6$, and $E_6$ singlets that arise in the string models, with the production and decay being produced via one-loop couplings to heavy vector-like matter states. Such vector-like matter states can provide viable dark matter candidates. This would be the case if the different 16 and 10 charges can be determined experimentally, they provide a distinct signatures of the heterotic-string models that utilise discrete Wilson lines to break the $E_6$ symmetry.

2 The string model

The string model was constructed in Ref. [24] and its details will not be repeated here. The construction of the model utilizes the free fermionic model building rules [3–5], and the notation that we use is prevalent in the literature (see e.g. [3–5,10–24] and references therein). The model is generated by using the classification methods developed in [25] for the classification of type IIB superstrings and extended in [26–30] for the classification of heterotic-string vacua with different SO(10) subgroups. The space of vacua is generated by working with a fixed set of boundary condition basis vectors and varying the GGSO coefficients [26–30], which are ±1 phases in the one-loop partition function. The $Z'$ model under consideration here was obtained in the class of Pati–Salam heterotic-string models, which are generated by a set of 13 boundary condition basis vectors $B = \{v_1, v_2, \ldots, v_{13}\}$. The basis vectors are shown in Eq. (2.1),

$$v_1 = 1 = \{\psi^\mu, \chi^{1,6}, y^{1,6}, \omega^{1,6}\},$$
$$v_2 = S = \{\psi^\mu, \chi^{1,6}\},$$
$$v_2+i = e_i = \{y^i, \omega^i|\bar{y}^i, \bar{\omega}^i\}, \ i = 1, \ldots, 6,$$
$$v_9 = z_1 = \{\bar{\phi}^{1,4}\},$$
$$v_{10} = z_2 = \{\bar{\phi}^{5,8}\},$$
$$v_{11} = b_1 = \{\chi^{34}, \chi^{56}, \gamma^{34}, \gamma^{56}\},$$
$$v_{12} = b_2 = \{\chi^{12}, \chi^{56}, \gamma^{12}, \gamma^{56}\},$$
$$v_{13} = \alpha = \{\bar{\psi}^{4,5}, \bar{\phi}^{1,2}\}.$$
In terms of the notation used in Eq. (2.2) the GGSO one-loop coefficients in the $N=1$ partition function are given by

$$c_{ij}^{(v_j|v_i)} = \exp[i\pi(v_j|v_i)].$$

The full massless spectrum of the model, and its charges under the four dimensional gauge group, is given in Ref. [24]. In Tables 1 and 2 we provide a glossary of the states in the model and their charges under the $SU(4) \times SU(2) \times U(1)_{\xi}$ gauge group. We remark that we changed the notation of [24] for the sextet fields from $D$ and $\bar{D}$ to $\Delta$ and $\bar{\Delta}$ to avoid confusion with the notation for the vector-like quarks below. We note that the sextet states are in the vector representation of $SO(6) \equiv SU(4)$. They are vector-like with respect to the Standard Model subgroup, but they are chiral with respect to $U(1)_{\xi}$.

The model is derived by fishing a self-dual model under the spinor–vector duality at the SO(10) level, i.e. prior to incorporation of the basis vector $\alpha$. The SO(10) model exhibits the self-duality property under the exchange of the total number of spinorial plus anti-spinorial, and the total number of vectorial SO(10) representations. This is in fact a key ingredient in the construction of the model and in the possibility of having an anomaly free $U(1)_{\xi}$ as part of a low scale $Z'$. The spinor–vector duality was observed in the classification of free fermionic vacua with SO(10) GUT group [41–44]. The statement is that for every vacuum with a total number #1 of twisted 16 $\oplus \overline{16}$ spinorial representations and a total number #2 twisted 10 vectorial representations, there exist another vacuum in which the two are interchanged. To understand the origin of this duality it is instrumental to consider the case in which SO(10) $\times U(1)_{\xi}$ is enhanced to $E_6$. The chiral and anti-chiral representations of $E_6$ are the 27 and 27 representations that decompose under SO(10) $\times U(1)_{\xi}$ as $27 = 16_{1/2} + 10_{-1} + 1_2$ and $\overline{27} = 16_{-1/2} + 10_1 + 1_{-2}$. Thus, in the case of $E_6$ the total number of 16 $\oplus$ $\overline{16}$ is equal to the total number of 10 representations, i.e. this case is self-dual under the spinor–vector duality map. In the case of $E_6$, $U(1)_{\xi}$ is anomaly free by virtue of its embedding in $E_6$, whereas in vacua with broken $E_6$, $U(1)_{\xi}$ is in general anomalous [45]. Furthermore, the case of $E_6$ correspond to a string vacuum with (2, 2) worldsheet supersymmetry. The $N=2$ worldsheet supersymmetry on the bosonic side of the heterotic string has a spectral flow operator that exchanges between the spinorial and vectorial components of the $E_6$ representations [43,44]. The vacua with broken $E_6$ symmetry only possess (2, 0) worldsheet supersymmetry. In these vacua the would be spectral flow operator induces the map between the spinor–vector dual vacua. The string vacua, however, also admit a class of self-dual vacua under the spinor–vector duality map without enhancement of the gauge symmetry to $E_6$. This is possible because the different spinorial

| Symbol | Fields in [24] | $SU(4) \times SU(2)_L \times SU(2)_R$ |
|--------|----------------|--------------------------------------|
| $H^+$  | $H_{12}^I$     | $(2, 2, 1, 1, 1)$                     |
| $H^-$  | $H_{12}^I$     | $(1, 1, 2, 2, 1)$                     |
| $H^+$  | $H_{13}^I$     | $(2, 2, 2, 1, 1)$                     |
| $H^-$  | $H_{13}^I$     | $(1, 1, 2, 2, 1)$                     |
| $H   $ | $H_{12}^I$     | $(2, 2, 1, 1, 1)$                     |
| $Z    $ | $Z_{i, i = 1, \ldots, 9}$ | $(1, 1, 8)$ |
and vectorial components that make up complete $E_6$ representations are obtained from different fixed points of the underlying $Z_2 \times Z_2$ orbifold. In this case the chiral spectrum still resides in complete $E_6$ representations, and $U(1)_\xi$ is anomaly free, but the gauge symmetry is not enhanced to $E_6$.

The Pati–Salam heterotic-string model generated by Eqs. (2.1) and (2.2) breaks the string matter states into the PS components, and as a result some states are projected out. However, as can be seen from Table 1 the twisted chiral matter spectrum of this model forms complete $E_6$ representations. It is noted that for $U(1)_\xi$ to be anomaly free only the chiral spectrum has to form complete 27 representations of $E_6$, whereas the string model may contain vector-like states that do not form complete $E_6$ representations. The observable and hidden gauge symmetries at the string scale are generated by untwisted states and are given by

observable: $SO(6) \times SO(4) \times U(1)_1 \times U(1)_2 \times U(1)_3$

hidden: $SO(4)^2 \times SO(8)$.

All the additional massless spacetime vector bosons that can enhance the observable and hidden gauge symmetries are projected out in this model due to the choice of GGSO phases in Eq. (2.2). The string model contains two anomalous $U(1)$s with

$$\text{Tr} U(1)_1 = 36 \quad \text{and} \quad \text{Tr} U(1)_3 = -36.$$  (2.3)

Consequently, the $E_6$ combination

$$U(1)_\xi = U(1)_1 + U(1)_2 + U(1)_3$$  (2.4)

is anomaly free and can be preserved as a component of an extra $Z'$ at lower scales. It should be emphasised that generally $U(1)_\xi$ is anomalous in the string models and therefore has to be broken in these models near the string scale [46,47]. It is anomaly free in the model generated by Eqs. (2.1) and (2.2) because the chiral spectrum forms complete 27 of $E_6$ multiplets.

As seen from Table 2 the model also contains vector-like states that transform under the hidden sector $SU(2)^4 \times SO(8)$ gauge group. They comprise four bi-doublets denoted by $H^\pm$, which carry $Q_\xi = \pm 1$ charges, 12 neutral bi-doublets, denoted by $H$, and five states that transform in the 8 representation of the hidden $SO(8)$ gauge group with $Q_\xi = 0$.

In the notation of Tables 1 and 2, the effective trilevel superpotential takes the form

$$W = w + w'.$$  (2.5)

where

$$w = \tilde{F} R F L h + \tilde{F} R F L \Delta + F L F _L \Delta + F R F R \Delta + F R \tilde{F} R \zeta + h h S + \Delta S + \tilde{D} D S + \Delta D \zeta + \zeta \zeta + S \tilde{\phi} \phi + S \phi \phi + S S \zeta.$$  (2.6)

$$w' = S H^0 H^0 + \tilde{S} H^0 H^0 + S H H + \phi H^0 H + \tilde{\phi} H^0 H + \phi \phi \zeta + \zeta H H + \zeta H^0 H^0 + Z Z \zeta.$$  (2.7)

where we have suppressed all generation and field indices.

As seen from Table 1 the string model contains the heavy Higgs states required to break the non-Abelian Pati–Salam symmetry [48]. These are $H = F_R$ and $\tilde{H}$, being a linear combination of the four $\tilde{F}_R$ fields. The decomposition of these fields in terms of the Standard Model group factors is given by

$$\tilde{H}(4, 1, 2) \to u_H \left(3, 1, \frac{2}{3}\right) + d_H \left(3, 1, -\frac{1}{3}\right) + \tilde{N}(1, 1, 0) + e_H (1, 1, -1).$$

The suppression of the left-handed neutrino masses favours the breaking of the Pati–Salam symmetry at the high scale. Schematically, the neutrino seesaw mass matrix in terms of the Standard Model components takes the generic form [16–18,49–51]

$$\begin{pmatrix} \nu \ N \ \zeta \end{pmatrix} \begin{pmatrix} 0 & (M_D) & 0 \\ 0 & 0 & \langle \tilde{N} \rangle \\ 0 & \langle \zeta \rangle & \langle \tilde{\zeta} \rangle \end{pmatrix} \begin{pmatrix} \nu \\ N \\ \zeta \end{pmatrix},$$  (2.8)

where generation indices are suppressed, and $M_D$ is the Dirac mass matrix arising from the couplings of the chiral fermions to the light Higgs bi-doublets. The underlying $SO(10)$ symmetry dictates that the Dirac mass term of the tau neutrino is proportional to that of the top quark [49–51]. Hence, adequate suppression of the tau neutrino mass favours high scale breaking of $SU(2)_R$. More intricate scenarios in which $SU(2)_R$ is broken at a lower scale, possibly even near the TeV scale, may be possible as well [52–54], however, for our purpose here we may assume that it is broken near the string scale. The breaking of the PS symmetry then leaves an unbroken $U(1)_Z'$ symmetry given by

$$U(1)_Z' = \frac{3}{10} U(1)_B - \frac{2}{5} U(1)_T + U(1)_\xi \notin SO(10),$$  (2.9)

which may remain unbroken down to low scales provided that $U(1)_\xi$ is anomaly free. Furthermore, cancellation of the anomalies mandates the existence of additional vector-like
quarks and leptons, arising from the vectorial 10 representation of SO(10), as well as the SO(10) singlet in the 27 of $E_6$. The spectrum below the PS breaking scale is displayed schematically in Table 3. The three right-handed neutrino $N^i_L$ fields are neutral under the gauge symmetry below the SU(2)$_R$ breaking scale and are decoupled from the low energy spectrum. This condition is specific to the extra $U(1)_{Z'}$ combination in Eq. (2.9). Here we assume that the spectrum is supersymmetric. We allowed for the possibility that the spectrum contains additional pairs of vector-like electroweak doublets and colour triplets. The spectrum may be compatible with GUT scale gauge coupling unification [55–57], where we may assume that the unification scale is either at the GUT or string scales [58], provided that there is an excess of one pair of vector-like electroweak doublets beyond the number of pairs of vector-like colour triplets.

This is possible in the free fermionic heterotic-string model, as noted from Table 1, is the existence of the exotic states $\phi_{1,2}$, and $\tilde{\phi}_{1,2}$. These states arise due to the breaking of $E_6$ by a discrete Wilson line in the string model. Such states do not arise in pure field theory GUT models, and may be a distinct signature of the specific string vacuum of Eqs. (2.1) and (2.2), i.e. they may be a distinct signature of the particular Wilson line used in this model. These exotic states are SO(10) singlets and therefore are neutral with respect to the Standard Model gauge group. The breaking of the $U(1)_{Z'}$ gauge symmetry may leave a discrete symmetry that forbids the decay of these exotic states to the lighter Standard Model states. This is the case if the $U(1)_{Z'}$ gauge symmetry is broken by one of the $S_i$ states. The mass scale of the exotic states, relative to the $Z'$ breaking scale, then determines the type and whether they can provide a viable dark matter candidate [59,60].

### Table 3 Spectrum and SU(3)$_C \times SU(2)_L \times U(1)_Y \times U(1)_{Z'}$ quantum numbers, with $i = 1, 2, 3$ for the three light generations. The charges are displayed in the normalisation used in free fermionic heterotic-string models

| Field | SU(3)$_C$ | SU(2)$_L$ | U(1)$_Y$ | U(1)$_{Z'}$ |
|-------|-----------|-----------|-----------|-------------|
| $Q^i_L$ | 3 | 2 | $+\frac{1}{6}$ | $-\frac{5}{6}$ |
| $\nu^i_L$ | 3 | 1 | $-\frac{2}{3}$ | $-\frac{1}{3}$ |
| $d^i_L$ | 3 | 1 | $+\frac{1}{3}$ | $-\frac{4}{3}$ |
| $e^i_L$ | 1 | 1 | $1$ | $-\frac{2}{3}$ |
| $L^i_L$ | 1 | 2 | $-\frac{1}{2}$ | $-\frac{1}{2}$ |
| $\Delta^i$ | 3 | 1 | $-\frac{1}{3}$ | $+\frac{5}{3}$ |
| $\tilde{\Delta}^i$ | 3 | 1 | $+\frac{1}{3}$ | $+\frac{5}{3}$ |
| $H_i$ | 1 | 2 | $-\frac{1}{2}$ | $+\frac{5}{2}$ |
| $\phi^i$ | 1 | 2 | $+\frac{1}{2}$ | $+\frac{5}{2}$ |
| $S^i$ | 1 | 1 | 0 | $-2$ |
| $h$ | 1 | 2 | $-\frac{1}{2}$ | $-\frac{4}{3}$ |
| $\tilde{h}$ | 1 | 2 | $+\frac{1}{2}$ | $+\frac{4}{3}$ |
| $\tilde{D}$ | 3 | 1 | $-\frac{1}{2}$ | $+\frac{5}{2}$ |
| $\bar{\tilde{D}}$ | 3 | 1 | $+\frac{1}{2}$ | $-\frac{4}{2}$ |
| $\phi$ | 1 | 1 | 0 | $-1$ |
| $\tilde{\phi}$ | 1 | 1 | 0 | $+1$ |
| $\zeta^i$ | 1 | 1 | 0 | 0 |

### 3 The di-photon events

The ATLAS [31] and CMS [32] experiments reported in December 2015 an excess in the production of di-photon events with a resonance around 750 GeV. It generated a flurry of activity with over 120 related papers since the announcement (for a partial list see e.g.: [64–116]). The statistical significance of the combined results is of the order of 3 sigma.

The more substantial indications are observed by the ATLAS experiment, which favours a rather broad width of the order of $\Gamma \sim 45$ GeV, which is not incompatible with the CMS results. The Landau-Yang theorem implies that only spin 0 or 2 resonance can decay into two photons. In the context of the perturbative heterotic-string construction the viable possibility is therefore a spin 0 resonance decaying into two photons. The production can be generated via gluon fusion, similar to the signal that led to the discovery of the Standard Model Higgs boson via $h \rightarrow 2\gamma$. There is no evidence at comparable energy scale of the resonance decaying into final states with any other particles, i.e. with $t\bar{t}, b\bar{b}, \ell^+\ell^-, ZZ, WW$, etc. A plausible explanation for the production and decay of such a resonance is via a Standard Model singlet scalar field that couples to heavy vector-like quark and lepton pairs. In is noted that indeed scenarios along this lines have been proposed for such a resonance [117–119].

Turning to the low energy spectrum of the model in Table 3 it is noted that the SO(10) singlets in the 27 of $E_6$ $S_i$, as well as the $E_6$ singlets $\zeta_i$, provide the needed fields. The low scale superpotential (2.6) gives rise to the terms

$$\lambda_{i,j,k}^D S^i D^j \bar{D}^k + \lambda_{i,j,k}^H S^i H^j H^k + M_{i} D^i D_i + M_{h} H_i H_i^i \, ,$$

(3.1)

$$\eta^D \zeta D \bar{D} + \eta^h \zeta h h + M_{D} D \bar{D} + \eta h h \, ,$$

(3.2)

where we allowed for the possibility that the couplings arise from terms that couple the vector-like states to the $S_i$ fields that carry $Q_{Z'} = -2$, as in Eq. (3.1), as well the coupling to the $\zeta_i$ fields that are neutral under $U(1)_{Z'}$. These terms can
generate the diphoton events via the diagram displayed in Fig. 1. Indeed, such terms are ubiquitous in the string models. The cubic level and higher terms in the superpotential are calculated by using the tools developed in [61–63]. In the model of Ref. [24] we find the terms for the states from Table 1 with couplings similar to those in Eq. (3.1),

\[ \Delta \Delta S + hhS = D_5 D_7 \chi_1^+ + D_3 D_4 \chi_1^+ + D_2 D_5 \chi_2^+ + D_4 D_7 \chi_2^+ + D_2 D_6 \chi_3^+ + D_1 D_7 \chi_5^+ + D_4 D_5 \chi_5^+ + D_1 D_2 \phi_{12} \\
+ D_2 D_3 \phi_{23} + D_1 D_3 \phi_{13} + D_4 D_4 \phi_{12} + D_5 D_5 \phi_{13} + D_6 D_6 \phi_{13} + D_7 D_7 \phi_{23} + h_2 h_2 \phi_{13} + h_3 h_3 \phi_{13} + h_1 h_1 \phi_{12} + h_1 h_2 \chi_5^+ , \]  

(3.3)

as well as couplings similar to those in Eq. (3.2),

\[ \Delta \Delta \zeta = + \bar{D}_1 D_7 \chi_5^- + \bar{D}_2 D_5 \chi_5^- + \bar{D}_2 D_6 \chi_3^- + \bar{D}_3 D_4 \chi_1^- + \bar{D}_2 \bar{D}_7 \phi_{12} + \bar{D}_1 \bar{D}_3 \phi_{13} + \bar{D}_2 \bar{D}_3 \phi_{23} + \bar{D}_3 \bar{D}_2 \phi_{23} + \bar{D}_1 \bar{D}_3 \phi_{13} + \bar{D}_3 \bar{D}_1 \phi_{13} + D_1 \bar{D}_6 \chi_4^- + D_6 \bar{D}_6 \xi_1 . \]  

(3.4)

The chiral spectrum of the model after SU(2)_R breaking gives rise to three copies of the chiral states shown in Table 3. The string model does not give rise, however, to the extra pair of vector-like Higgs doublets, which are instrumental for gauge coupling unification. The reason is that our string model uses symmetric boundary conditions rather than asymmetric boundary conditions [53,54].

The VEVs of the heavy Higgs fields \( \mathcal{H}, \bar{\mathcal{H}} \), the weak hypercharge combination given by

\[ U(1)_Y = \frac{1}{2} U(1)_B - L + U(1)_{T_3_R} , \]  

(3.5)

which is unbroken, as well as the \( U(1)_{Z'} \) combination given in Eq. (2.9), which is orthogonal to the weak hypercharge combination. The scalar component of one of the \( S_i \) fields provide the Higgs field that breaks the extra \( U(1) \) symmetry \( \langle S \rangle = v' \). Provided that this VEV is of the order of the TeV scale, say \( v' = 5 \) TeV, then it ensures that the \( U(1)_{Z'} \) remains unbroken down to the TeV scale. Furthermore, ensuring that the extra \( U(1)_{Z'} \) is anomaly free mandates that the extra vector-like quarks and leptons obtain their mass of the order of \( O(v') \) from the couplings in Eq. (3.1). Thus, all the ingredients needed to generate the characteristics of the diphoton events naturally exist in the string model. The model then associates the diphoton events with the existence of additional \( U(1)_{Z'} \) gauge symmetry at \( O(v') \). With \( M_{Z'} \sim g' v' \) and \( M_S \sim \lambda' v' \) being naively the masses of the heavy \( Z' \) vector boson and the Higgs field \( S \), respectively, we see that in this model the masses of the \( Z' \) vector field and the \( S \) scalar field are closely related. These characteristics fit well with both a di-boson excess at 1.9 TeV [33–35] as well as with the di-photon excess at 750 GeV [31,32]. Furthermore, the model predicts the existence of the additional vector-like quarks and leptons in the same vicinity. The existence of the \( U(1)_{Z'} \) symmetry at the TeV scale, and the associated anomaly cancellation requirement, naturally explain the existence of the vector-like quarks and leptons at the \( U(1)_{Z'} \) breaking scale. However, as follows from Eqs. (3.2) and (3.4), the diphoton events can also be mediated in the string models by scalars with \( Q_{Z'} = 0 \), and in this case the mass scale of the vector-like states is disassociated from the \( U(1)_{Z'} \) breaking scale. This is a less appealing scenario, but one which is allowed in the string model. In this case \( U(1)_{Z'} \) may be broken at a high scale along a flat direction by utilising the VEVs for, say \( \tilde{\chi}_4^+ \) and \( \chi_S^- \).
4 Conclusions

Early indications from LHC run 1 and 2 provide evidence for excess in the di-photon channel with a resonance of the order of 750 GeV, and a less significant di-boson excess at 1.9 TeV. Such signals fit naturally in our heterotic-string derived model, with a high scale seesaw mechanism and a low scale $U(1)_{Z'}$ breaking. Furthermore, the model predicts the existence of additional vector-like quarks and leptons related to the $U(1)_{Z'}$ breaking scale.

Some of the proposals to explain the di-photon excess are in line with the explanation employed in our paper (see e.g. [117–119]), and alternative scenarios have been proposed as well (see e.g. [64–116]). Some of these alternative scenarios employ a composite scalar singlet. The string derived model [24] does contain hidden sector fields, charged under $U(1)_R$, which can form composites that can mimic the charges of the singlet field $S_1$. Investigating whether this can provide an alternative scenario in the model at hand is left for future work.

The existence of vector-like quarks at the TeV scale poses a challenge when confronted with proton decay limits. Generically we anticipate that these states couple to the light quarks and may generate proton decay. How to avoid this conundrum remains a puzzle. Some plausible suggestions include the existence of local discrete symmetries that forbid the omnious couplings [120], and the special placement of matter fields in unified multiplets [121].

Alternative explanations have also been proposed in the case of models with a low string scale [36–40]. The low scale string scenarios give rise to additional Kaluza–Klein states at 1.9 TeV. Such signals fit naturally in our heterotic-string derived model, with a high scale seesaw mechanism and a low scale $U(1)_{Z'}$ breaking. Furthermore, the model predicts the existence of additional vector-like quarks and leptons related to the $U(1)_{Z'}$ breaking scale.

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Alternative explanations have also been proposed in the case of models with a low string scale [36–40]. The low scale string scenarios give rise to additional Kaluza–Klein and heavy string modes and therefore will easily be discerned from the heterotic-string scenarios. Explorations into the multi-TeV regime will adjudicate between the competing scenarios. We await the return of the collider.

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