LTH–1073

The 750 GeV di–photon LHC excess and Extra Z's in Heterotic–String Derived Models

Alon E. Faraggi[†1] and John Rizos[‡2]

† *Dept. of Mathematical Sciences, University of Liverpool, Liverpool L69 7ZL, UK*

‡ *Department of Physics, University of Ioannina, GR45110 Ioannina, Greece*

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 Z' models, where the di–photon excess may be connected with the Standard Model singlet scalar responsible for the Z' symmetry breaking, whereas the di–boson excess arises from production of the extra vector boson. Additional vector–like states in the string 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 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 $U(1)_{Z'}$ breaking.

[1] E-mail address: alon.faraggi@liv.ac.uk
[2] E-mail address: irizos@uoi.gr

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 [1, 2]. The free fermionic formulation [3] 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 $Z_2 \times Z_2$ 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 [4]. The viable models constructed to date include the flipped $SU(5)$ (FSU5) [5]; the standard–like models (SLM) [6]; the Pati–Salam models (PS) [7]; and the left–right symmetric models (LRS) [8]; whereas the $SU(4) \times SU(2) \times U(1)$ subgroup did not produce viable models [9]. All these models give rise to extra observable gauge symmetries at the string scale. Flavour universal symmetries typically arise from the $SO(10)$ and E_6 group factors. However, preserving an unbroken extra gauge symmetry down to low scales has proven to be elusive in the string models. The reasons being that suppression of left–handed neutrino masses favours the breaking of lepton number at a high scale, whereas the $U(1)_\zeta$ gauge symmetry in the decomposition of $E_6 \to SO(10) \times U(1)_\zeta$ tends to be anomalous in the string models and therefore cannot remain unbroken down to low scales. This extra $U(1)$ in the string models is typically anomalous because of the projection of some states from the spectrum by the Generalised GSO (GGSO) projections, *i.e.* anomaly cancellation requires that the chiral spectrum appears in complete E_6 representations. However, the breaking of E_6 at the string scale mandates that the chiral states exist in incomplete E_6 multiplets.

Recently, however, we were able to construct a heterotic–string model in which the desired symmetry is anomaly free [10]. The derivation of this model utilises the spinor–vector duality that was discovered in free fermionic heterotic–string models, and was obtained by using the classification methodology developed in [11–14]. The model of ref. [10] is obtained from a self–dual $SO(10)$ model under the exchange of the total number of spinorial $16 \oplus \overline{16}$ and vectorial 10 representations of $SO(10)$. This is the condition that one has if the $SO(10) \times U(1)_\zeta$ symmetry is enhanced to E_6. However, in the model of ref. [10] the $SO(10)$ symmetry is not enhanced to E_6. This is the case in the free fermionic model if the different 16 and 10+1 representations, that would form a complete 27 representation of E_6, are obtained from different fixed points of the underlying $Z_2 \times Z_2$ orbifold [10]. 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 $U(1)_\zeta$ anomaly free. We remark that while complete E_6 multiplets is sufficient for $U(1)_\zeta$ to be anomaly free, it might not be necessary and alternative possibilities may exist.

The ATLAS and CMS collaborations [15, 16] 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 are precisely those required from anomaly cancellation of the extra Z' gauge symmetry. Indications for an extra Z' of order 2 TeV have been earlier suggested by the ATLAS and CMS experiments [17]. They have not been substantiated by the run II data, but the possibility of an extra Z' at that scale nevertheless remains. The signature of the string model with a low scale Z' in our model is characterised by the scalar resonance, by the extra Z' and by the additional vector–like matter at the multi–TeV scale, required by anomaly cancellation.

Alternative string constructions that may account for such excesses have been recently suggested [18], based on low scale string models and F–theory scenarios. Low scale string models give rise to Kaluza–Klein and heavy string states. Therefore, future colliders will be able to discriminate between perturbative heterotic–string models and low scale string scenarios. Additionally, heterotic–string models give rise to states that do not satisfy the E_6 quantisation of the Z' charges. Such states arise in string models due to the breaking of the E_6 symmetry by Wilson lines and can produce viable dark matter candidates. This would be the case if the Z' symmetry is broken by a state that carries standard E_6 charge. A residual discrete symmetry then forbids the decay of the exotic string state to states that carry the standard E_6 charges. Such exotic states at the multi–TeV scale can provide viable thermal relics. On the other hand, if their exotic Z' 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. [10] and its details will not be repeated here. The construction of the model utilizes the free fermionic model building rules [3], and the notation that we use is prevalent in the literature (see *e.g.* [3, 5–10] and references therein). The model is generated by using the classification methods developed in [11] for the classification of type IIB superstrings and extended in [12–14] 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 [12–14], 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 thirteen boundary condition basis vectors $B = \{v_1, v_2, \ldots, v_{13}\}$. The basis vectors are shown in eq. (2.1),

$$
\begin{aligned}
v_1 = 1 &= \{\psi^\mu,\ \chi^{1,\dots,6}, y^{1,\dots,6}, \omega^{1,\dots,6} | \\
&\qquad \bar{y}^{1,\dots,6}, \bar{\omega}^{1,\dots,6}, \bar{\eta}^{1,2,3}, \bar{\psi}^{1,\dots,5}, \bar{\phi}^{1,\dots,8}\}, \\
v_2 = S &= \{\psi^\mu, \chi^{1,\dots,6}\}, \\
v_{2+i} = e_i &= \{y^i, \omega^i | \bar{y}^i, \bar{\omega}^i\},\ i = 1, \dots, 6, \\
v_9 = z_1 &= \{\bar{\phi}^{1,\dots,4}\}, \\
v_{10} = z_2 &= \{\bar{\phi}^{5,\dots,8}\}, \\
v_{11} = b_1 &= \{\chi^{34}, \chi^{56}, y^{34}, y^{56} | \bar{y}^{34}, \bar{y}^{56}, \bar{\eta}^1, \bar{\psi}^{1,\dots,5}\}, \\
v_{12} = b_2 &= \{\chi^{12}, \chi^{56}, y^{12}, y^{56} | \bar{y}^{12}, \bar{y}^{56}, \bar{\eta}^2, \bar{\psi}^{1,\dots,5}\}, \\
v_{13} = \alpha &= \{\bar{\psi}^{4,5}, \bar{\phi}^{1,2}\}.
\end{aligned}
\tag{2.1}
$$

The fermions appearing in the curly brackets in eq. (2.1) are periodic, whereas those that do not appear are antiperiodic. The first twelve basis vectors in (2.1) generate the space of vacua with unbroken $SO(10)$ symmetry [12], whereas the thirteenth basis vector breaks the $SO(10)$ symmetry to the Pati–Salam subgroup [13]. The one–loop GGSO phases between the basis vectors are given by a 13×13 matrix. Only the terms above the diagonal are independent, whereas those on the diagonal and below are fixed by modular transformations [3]. Additional constraints, such as requiring the vacuum to possess space–time supersymmetry, fixes additional phases and leaves a total of 66 independent phases. Using a random generation algorithm we can generate random choices of the independent phases. By imposing some physical criteria on the desired model, we can fish out models with desired physical characteristics. These include the absence of symmetry enhancing spacetime vector bosons and exotic fractionally charged states from the massless spectrum. A choice of GGSO phases that produces these desired results is given by:

$$
(v_i|v_j) \;=\; \begin{array}{c} \\ 1 \\ S \\ e_1 \\ e_2 \\ e_3 \\ e_4 \\ e_5 \\ e_6 \\ b_1 \\ b_2 \\ z_1 \\ z_2 \\ \alpha \end{array}
\begin{array}{c}
\begin{array}{ccccccccccccc} 1 & S & e_1 & e_2 & e_3 & e_4 & e_5 & e_6 & b_1 & b_2 & z_1 & z_2 & \alpha \end{array} \\
\left(\begin{array}{ccccccccccccc}
1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 \\
1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 \\
1 & 1 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 1 \\
1 & 1 & 0 & 0 & 0 & 0 & 0 & 1 & 0 & 0 & 0 & 1 & 0 \\
1 & 1 & 0 & 0 & 0 & 1 & 0 & 0 & 0 & 0 & 0 & 1 & 1 \\
1 & 1 & 0 & 0 & 1 & 0 & 0 & 0 & 0 & 0 & 1 & 0 & 0 \\
1 & 1 & 0 & 0 & 0 & 0 & 0 & 1 & 0 & 0 & 0 & 1 & 1 \\
1 & 1 & 0 & 1 & 0 & 0 & 1 & 0 & 0 & 0 & 1 & 0 & 0 \\
1 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 1 & 1 & 0 & 0 & 0 \\
1 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 1 & 1 & 0 & 0 & 1 \\
1 & 1 & 0 & 0 & 0 & 1 & 0 & 1 & 0 & 0 & 1 & 1 & 0 \\
1 & 1 & 0 & 1 & 1 & 0 & 1 & 0 & 0 & 0 & 1 & 1 & 0 \\
1 & 1 & 1 & 0 & 1 & 0 & 1 & 0 & 1 & 0 & 1 & 0 & 1
\end{array}\right)
\end{array}
\tag{2.2}
$$

In terms of the notation used in eq. (2.2) the GGSO one–loop coefficients in the $N = 1$ partition function are given by

$$c\begin{bmatrix} vi \\ v_j \end{bmatrix} \; = \; \exp[i\pi(v_i|v_j)].$$

The full massless spectrum of the model, and its charges under the four dimensional gauge group, is given in ref [10]. In tables 1 and 2 we provide a glossary of the states in the model and their charges under the $SU(4) \times SO(4) \times U(1)_\zeta$ gauge group. We remark that we changed the notation of [10] for the sextet fields from D and $\bar{D}$ to Δ 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 are chiral with respect to $U(1)_\zeta$.

Symbol	Fields in [10]	$SU(4) \times SU(2)_L \times SU(2)_R$	$U(1)_\zeta$
F_L	F_{1L}, F_{2L}, F_{3L}	$(\mathbf{4}, \mathbf{2}, \mathbf{1})$	$+\frac{1}{2}$
F_R	F_{1R}	$(\mathbf{4}, \mathbf{1}, \mathbf{2})$	$-\frac{1}{2}$
$\bar{F}_R$	$\bar{F}_{1R}, \bar{F}_{2R}, \bar{F}_{3R}, \bar{F}_{4R}$	$(\bar{\mathbf{4}}, \mathbf{1}, \mathbf{2})$	$+\frac{1}{2}$
h	h_1, h_2, h_3	$(\mathbf{1}, \mathbf{2}, \mathbf{2})$	-1
Δ	$D_1, \ldots, D_7$	$(\mathbf{6}, \mathbf{1}, \mathbf{1})$	-1
$\bar{\Delta}$	$\bar{D}_1, \bar{D}_2, \bar{D}_3, \bar{D}_6$	$(\mathbf{6}, \mathbf{1}, \mathbf{1})$	$+1$
S	$\Phi_{12}, \Phi_{13}, \Phi_{23}, \chi^+_1, \chi^+_2, \chi^+_3, \chi^+_5$	$(\mathbf{1}, \mathbf{1}, \mathbf{1})$	$+2$
$\bar{S}$	$\bar{\Phi}_{12}, \bar{\Phi}_{13}, \bar{\Phi}_{23}, \bar{\chi}^+_4$	$(\mathbf{1}, \mathbf{1}, \mathbf{1})$	-2
ϕ	ϕ_1, ϕ_2	$(\mathbf{1}, \mathbf{1}, \mathbf{1})$	$+1$
$\bar{\phi}$	$\bar{\phi}_1, \bar{\phi}_2$	$(\mathbf{1}, \mathbf{1}, \mathbf{1})$	-1
ζ	$\Phi^-_{12}, \Phi^-_{13}, \Phi^-_{23}, \bar{\Phi}^-_{12}, \bar{\Phi}^-_{13}, \bar{\Phi}^-_{23}$	$(\mathbf{1}, \mathbf{1}, \mathbf{1})$	0
	$\chi^-_1, \chi^-_2, \chi^-_3, \bar{\chi}^-_4, \chi^-_5$		
	$\zeta_i, \bar{\zeta}_i, i = 1, \ldots, 9$		
	$\Phi_i, i = 1, \ldots, 6$		

Table 1: Observable sector field notation and associated states in [10].

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 α. 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 $U(1)_\zeta$ 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 [19,20]. 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.

Symbol	Fields in [10]	$SU(2)^4 \times SO(8)$	$U(1)_\zeta$
H^+	H^3_{12}	$(\mathbf{2},\mathbf{2},\mathbf{1},\mathbf{1},\mathbf{1})$	$+1$
	H^2_{34}	$(\mathbf{1},\mathbf{1},\mathbf{2},\mathbf{2},\mathbf{1})$	$+1$
H^-	H^2_{12}	$(\mathbf{2},\mathbf{2},\mathbf{1},\mathbf{1},\mathbf{1})$	-1
	H^3_{34}	$(\mathbf{1},\mathbf{1},\mathbf{2},\mathbf{2},\mathbf{1})$	-1
H	H^1_{12}	$(\mathbf{2},\mathbf{2},\mathbf{1},\mathbf{1},\mathbf{1})$	0
	$H^i_{13}, i=1,2,3$	$(\mathbf{2},\mathbf{1},\mathbf{2},\mathbf{1},\mathbf{1})$	0
	$H^i_{14}, i=1,2,3$	$(\mathbf{2},\mathbf{1},\mathbf{1},\mathbf{2},\mathbf{1})$	0
	H^1_{23}	$(\mathbf{1},\mathbf{2},\mathbf{2},\mathbf{1},\mathbf{1})$	0
	H^1_{24}	$(\mathbf{1},\mathbf{2},\mathbf{1},\mathbf{2},\mathbf{1})$	0
	$H^i_{34}, i=1,4,5$	$(\mathbf{1},\mathbf{1},\mathbf{2},\mathbf{2},\mathbf{1})$	0
Z	$Z_i, i=1,\ldots,$	$(\mathbf{1},\mathbf{1},\mathbf{8})$	0

Table 2: Hidden sector field notation and associated states in [10].

To understand the origin of this duality it is instrumental to consider the case in which $SO(10)\times U(1)_\zeta$ is enhanced to E_6. The chiral and anti–chiral representations of E_6 are the 27 and $\overline{27}$ representations that decompose under $SO(10)\times U(1)_\zeta$ as $27 = 16_{1/2} + 10_{-1} + 1_2$ and $\overline{27} = \overline{16}_{-1/2} + 10_1 + 1_{-2}$. 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, $U(1)_\zeta$ is anomaly free by virtue of its embedding in E_6, whereas in vacua with broken E_6, $U(1)_\zeta$ is in general anomalous [21]. 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 [20]. 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 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 $U(1)_\zeta$ is anomaly free, but the gauge symmetry is not enhanced to E_6.

The Pati–Salam heterotic string model generated by eqs. (2.1,2.2) 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 $U(1)_\zeta$ 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:

$$
\begin{aligned}
\text{observable} &: \quad SO(6)\times SO(4)\times U(1)_1\times U(1)_2\times U(1)_3 \\
\text{hidden} &: \quad SO(4)^2\times SO(8)
\end{aligned}
$$

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

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

Consequently, the E_6 combination

$$
U(1)_\zeta = U(1)_1 + U(1)_2 + U(1)_3 \tag{2.4}
$$

is anomaly free and can be preserved as a component of an extra Z' at lower scales. It should be emphasised that generically $U(1)_\zeta$ is anomalous in the string models and therefore has to be broken in these models near the string scale [22]. It is anomaly free in the model generated by eqs. (2.1,2.2) 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 bidoublets denoted by $H^\pm$ that carry $Q_\zeta = \pm 1$ charges, 12 neutral bidoublets denoted by H and five states that transform in the 8 representation of the hidden $SO(8)$ gauge group with $Q_\zeta = 0$.

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

$$
W = w + w' \tag{2.5}
$$

where

$$
\begin{aligned}
w &= \bar{F}_R F_L h + \bar{F}_R \bar{F}_R \Delta + F_L F_L \Delta + F_R F_R \bar{\Delta} + F_R \bar{F}_R \zeta + hhS \\
&\quad + \Delta\Delta S + \bar{D}\bar{D}\bar{S} + \Delta\bar{D}\zeta + \zeta\zeta\zeta + S\bar{\phi}\bar{\phi} + \bar{S}\phi\phi + S\bar{S}\zeta \ , \qquad (2.6) \\
w' &= SH^-H^+ + \bar{S}H^+H^+ + SHH + \phi H^- H + \bar{\phi}H^+H + \phi\bar{\phi}\zeta \\
&\quad + \zeta HH + \zeta H^+H^- + ZZ\zeta \ , \qquad (2.7)
\end{aligned}
$$

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 [23]. These are $\mathcal{H} = F_R$ and $\bar{\mathcal{H}}$, being a linear combination of the four $\bar{F}_R$ fields. The decomposition of these fields in terms of the Standard Model group factors is given by:

$$\begin{aligned}
\bar{\mathcal{H}}(\bar{\mathbf{4}},\mathbf{1},\mathbf{2}) &\rightarrow u^c_H\left(\bar{\mathbf{3}},\mathbf{1},\frac{2}{3}\right)+d^c_H\left(\bar{\mathbf{3}},\mathbf{1},-\frac{1}{3}\right)+\bar{\mathcal{N}}\left(\mathbf{1},\mathbf{1},0\right)+e^c_H\left(\mathbf{1},\mathbf{1},-1\right)\\
\mathcal{H}\left(\mathbf{4},\mathbf{1},\mathbf{2}\right) &\rightarrow u_H\left(\mathbf{3},\mathbf{1},-\frac{2}{3}\right)+d_H\left(\mathbf{3},\mathbf{1},\frac{1}{3}\right)+\mathcal{N}\left(\mathbf{1},\mathbf{1},0\right)+e_H\left(\mathbf{1},\mathbf{1},1\right)
\end{aligned}$$

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 [7, 24]

$$\begin{pmatrix}\nu, & N, & \zeta\end{pmatrix}\begin{pmatrix}0 & (M_{_D}) & 0\\ (M_{_D}) & 0 & \langle\bar{\mathcal{N}}\rangle\\ 0 & \langle\bar{\mathcal{N}}\rangle & \langle\zeta\rangle\end{pmatrix}\begin{pmatrix}\nu\\ N\\ \zeta\end{pmatrix}, \tag{2.8}$$

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 [24]. 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 [25], 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

$$U(1)_{Z'} \;=\; \frac{3}{10}U(1)_{B-L}-\frac{2}{5}U(1)_{T_{3_R}}-U(1)_\zeta \;\notin\; SO(10), \tag{2.9}$$

that may remain unbroken down to low scales provided that $U(1)_\zeta$ 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 Higgs doublets and colour triplets. The spectrum may be compatible with GUT scale gauge coupling unification [27], where we may assume that the unification scale is either at the GUT or string scales [28], 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 models due to the doublet–triplet splitting mechanism that operates in the string models [26]. Additionally, we allowed for the possibility of light states that are neutral under the $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{Z'}$ low scale gauge group. The $U(1)_{Z'}$ gauge symmetry can be broken at low scales by the VEV of the $SO(10)$ singlets S_i and/or $\phi_{1,2}$.

Field	$SU(3)_C$	$\times SU(2)_L$	$U(1)_Y$	$U(1)_{Z'}$
Q_L^i	3	2	$+\frac{1}{6}$	$-\frac{2}{5}$
u_L^i	$\bar{3}$	1	$-\frac{2}{3}$	$-\frac{2}{5}$
d_L^i	$\bar{3}$	1	$+\frac{1}{3}$	$-\frac{4}{5}$
e_L^i	1	1	$+1$	$-\frac{2}{5}$
L_L^i	1	2	$-\frac{1}{2}$	$-\frac{4}{5}$
D^i	3	1	$-\frac{1}{3}$	$+\frac{4}{5}$
$\bar{D}^i$	$\bar{3}$	1	$+\frac{1}{3}$	$+\frac{6}{5}$
H^i	1	2	$-\frac{1}{2}$	$+\frac{6}{5}$
$\bar{H}^i$	1	2	$+\frac{1}{2}$	$+\frac{4}{5}$
S^i	1	1	0	-2
h	1	2	$-\frac{1}{2}$	$-\frac{4}{5}$
$\bar{h}$	1	2	$+\frac{1}{2}$	$+\frac{4}{5}$
$\mathcal{D}$	3	1	$-\frac{1}{3}$	$+\frac{4}{5}$
$\bar{\mathcal{D}}$	$\bar{3}$	1	$+\frac{1}{3}$	$-\frac{4}{5}$
ϕ	1	1	0	-1
$\bar{\phi}$	1	1	0	$+1$
ζ^i	1	1	0	0

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.*

A distinct property of the free fermionic heterotic–string model, as noted from table 1, is the existence of the exotic states $\phi_{1,2}$, and $\bar{\phi}_{1,2}$. 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, 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 [29].

3 The di–photon events

The ATLAS [15] and CMS [16] experiments reported in December 2015 an excess in the 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 [31]. 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 $h \to 2\gamma$. There is no evidence at comparable energy scale of the resonance decaying into final states with any other particles, *i.e.* with $t\bar{t}$, $b\bar{b}$, $\ell\bar{\ell}$, 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 [32].

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 ζ_i, provide the needed fields. The low scale superpotential 2.6 gives rise to the terms

$$\lambda^D_{ijk} S^i D^j \bar{D}^k + \lambda^H_{ijk} S_i H^j \bar{H}^k + M^D_i D_i \bar{D}_i + M_H H_i \bar{H}_i \ , \tag{3.1}$$

$$\eta^{\mathcal{D}} \zeta \mathcal{D} \bar{\mathcal{D}} + \eta^h \zeta h \bar{h} + M_{\mathcal{D}} \mathcal{D} \bar{\mathcal{D}} + M_h h \bar{h} \ , \tag{3.2}$$

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 ζ_i fields that are neutral under $U(1)_{Z'}$. These terms can generate the diphoton events via the diagram displayed in figure 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 [30]. In the model of ref. [10] we find the terms for the states from table 1 with couplings similar to those

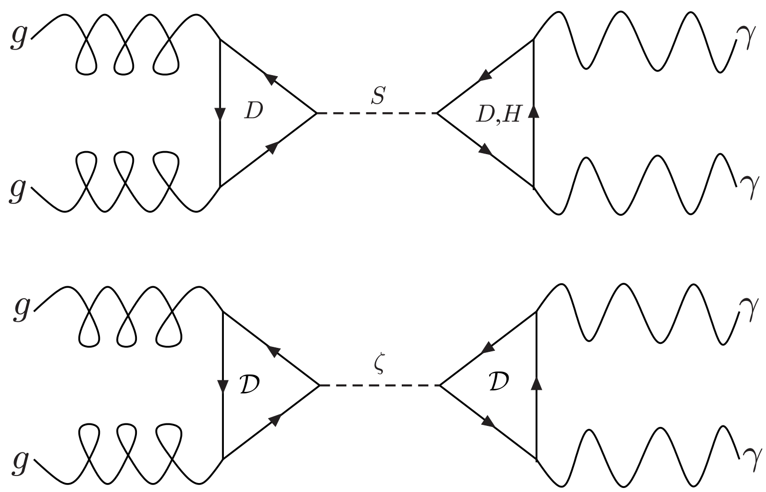

Figure 1: Production and diphoton decay of the $SO(10)$ singlet state. The top diagram corresponds to production via the terms in eq. (3.1) in the case with a low scale $U(1)_{Z'}$, whereas the bottom diagram corresponds to production via the terms in eq. (3.2).

in eq. (3.1)

$$\begin{aligned}
\Delta\Delta S + hhS \quad = \quad & D_5\, D_7\, \chi_1^+ + D_3\, D_4\, \chi_1^+ + D_2\, D_5\, \chi_2^+ + D_4\, D_7\, \chi_2^+ + D_2\, D_6\, \chi_3^+ \\
& + D_1\, D_7\, \chi_5^+ + D_4\, D_5\, \chi_5^+ + D_1\, D_2\, \Phi_{12} + D_2\, D_3\, \Phi_{23} + D_1\, D_3\, \Phi_{13} \\
& + D_4\, D_4\, \Phi_{12} + D_5\, D_5\, \Phi_{13} + D_6\, D_6\, \Phi_{13} + D_7\, D_7\, \Phi_{23} \\
& + h_2\, h_2\, \Phi_{13} + h_3\, h_3\, \Phi_{13} + h_1\, h_1\, \Phi_{12} + h_1\, h_2\, \chi_5^+ \qquad (3.3)
\end{aligned}$$

as well as couplings similar to those in eq. (3.2)

$$\begin{aligned}
\Delta\bar{\Delta}\zeta \quad = \quad & +\bar{D}_1\, D_7\, \chi_5^- + \bar{D}_2\, D_5\, \chi_2^- + \bar{D}_2\, D_6\, \chi_3^- + \bar{D}_3\, D_4\, \chi_1^- \\
& + D_2\, \bar{D}_1\, \bar{\Phi}_{12}^- + D_1\, \bar{D}_2\, \bar{\Phi}_{12}^- + D_2\, \bar{D}_3\, \Phi_{23}^- + D_3\, \bar{D}_2\, \bar{\Phi}_{23}^- \\
& + D_1\, \bar{D}_3\, \Phi_{13}^- + D_3\, \bar{D}_1\, \bar{\Phi}_{13}^- + D_1\, \bar{D}_6\, \bar{\chi}_4^- + D_6\, \bar{D}_6\, \zeta_1 \qquad (3.4)
\end{aligned}$$

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 being that our string model uses symmetric boundary conditions rather than asymmetric boundary conditions [26].

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

$$U(1)_Y = \frac{1}{2}U(1)_{B-L} + U(1)_{T_{3_R}} \qquad (3.5)$$

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 the Higgs field that breaks the extra $U(1)$ symmetry $\langle \mathcal{S} \rangle = v'$. Provided that this VEV is of the order of the TeV scale, say $v' = 5$ TeV, then 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 $M_{\mathcal{S}} \sim \lambda' v'$ being naively the masses of the heavy Z' vector boson and the Higgs field $\mathcal{S}$, respectively, we have that in this model the masses of the Z' vector field and the $\mathcal{S}$ scalar field are closely related. These characteristics fit well with both a di–boson excess at 1.9 TeV [17] as well as with the di–photon excess at 750 GeV [15, 16]. 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. (3.2,3.4), 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 $\bar{\chi}^+_4$ and χ^+_5.

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.* [32]), and alternative scenarios have been proposed as well (see *e.g.* [31]), Some of these alternative scenarios employ a composite scalar singlet. The string derived model [10] does contain hidden sector fields, charged under $U(1)_\zeta$, that can form composites that can mimic the charges of the singlet field S_i. 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 ominous couplings [33], and the special placement of matter fields in unified multiplets [34].

Alternative explanations have also been proposed in the case of models with a low string scale [18]. The low scale string scenarios give rise to additional Kaluza–Klein and heavy string modes and therefore will be easily 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.

Acknowledgments

AEF thanks Shmuel Nussinov for discussions and the theoretical physics groups at Oxford University, the Weizmann Institute and Tel Aviv University for hospitality. AEF is supported in part by the STFC (ST/L000431/1).

References

[1] D.J. Gross, J.A. Harvey, E.J. Martinec and R. Rohm, *Nucl. Phys.* **B267** (1986) 75.

[2] P. Candelas, G.T. Horowitz, A. Strominger and E. Witten, *Nucl. Phys.* **B258** (1985) 46.

[3] H. Kawai, D.C. Lewellen, and S.H.-H. Tye, *Nucl. Phys.* **B288** (1987) 1;
I. Antoniadis, C. Bachas, and C. Kounnas, *Nucl. Phys.* **B289** (1987) 87;
I. Antoniadis and C. Bachas, *Nucl. Phys.* **B289** (1987) 87.

[4] A.E. Faraggi, *Phys. Lett.* **B326** (1994) 62; *Phys. Lett.* **B544** (2002) 207;
E. Kiritsis and C. Kounnas, *Nucl. Phys.* **B503** (1997) 117;
A.E. Faraggi, S. Forste and C. Timirgaziu, JHEP **0**608, (2006) 057.

[5] I. Antoniadis, J. Ellis, J. Hagelin and D.V. Nanopoulos *Phys. Lett.* **B231** (1989) 65.

[6] A.E. Faraggi, D.V. Nanopoulos and K. Yuan, *Nucl. Phys.* **B335** (1990) 347;
A.E. Faraggi, *Phys. Lett.* **B278** (1992) 131; *Nucl. Phys.* **B387** (1992) 239;
G.B. Cleaver, A.E. Faraggi and D.V. Nanopoulos, *Phys. Lett.* **B455** (1999) 135;
A.E. Faraggi, E. Manno and C. Timirgaziu, *Eur. Phys. Jour.* **C50** (2007) 701.

[7] I. Antoniadis, G.K. Leontaris and J. Rizos, *Phys. Lett.* **B245** (1990) 161;
G.K. Leontaris and J. Rizos, *Nucl. Phys.* **B554** (1999) 3;
K. Christodoulides, A.E. Faraggi and J. Rizos, *Phys. Lett.* **B702** (2011) 81.

[8] G.B. Cleaver, A.E. Faraggi and C. Savage, *Phys. Rev.* **D63** (2001) 066001;
G.B. Cleaver, D.J. Clements and A.E. Faraggi *Phys. Rev.* **D65** (2002) 106003.

[9] G.B. Cleaver, A.E. Faraggi and S.E.M. Nooij, *Nucl. Phys.* **B672** (2003) 64;
A.E. Faraggi and H. Sonmez, *Phys. Rev.* **D91** (2015) 066006;
H. Sonmez, arXiv:150301193.

[10] A.E. Faraggi and J. Rizos, *Nucl. Phys.* **B895** (2015) 233.

[11] A. Gregori, C. Kounnas and J. Rizos, *Nucl. Phys.* **B549** (1999) 16.

[12] A.E. Faraggi, C. Kounnas, S.E.M. Nooij and J. Rizos, *Nucl. Phys.* **B695** (2004) 41;
A.E. Faraggi, C. Kounnas and J. Rizos, *Phys. Lett.* **B648** (2007) 84.

[13] B. Assel, K. Christodoulides, A.E. Faraggi, C. Kounnas and J. Rizos, *Phys. Lett.* **B683** (2010) 306; *Nucl. Phys.* **B844** (2011) 365.

[14] A.E. Faraggi, J. Rizos and H. Sonmez, *Nucl. Phys.* **B886** (2014) 202.

[15] ATLAS Collaboration, G. Aad *et al*, *ATLAS–CONF–2015–081.*

[16] CMS Collaboration, S. Chatrchyan *et al*, *CMS PAS EXO–15–004.*

[17] G. Aad *et al.* [ATLAS Collaboration], arXiv:1506.00962;
V. Khachatryan *et al.* [CMS Collaboration], JHEP **1**408, (2014) 173;
V. Khachatryan *et al.* [CMS Collaboration], JHEP **1**408, (2014) 174.

[18] J.J. Heckman, arXiv:1512.06773;
L.A. Anchordoqui *et al*, arXiv:1512.08502;
L.E. Ibanez and V. Martin-Lozano, arXiv:1512.08777;
M. Cvetic, J. Halverson and P. Langacker, arXiv:1512.07622;
E. Palti, arXiv:1601.00285.

[19] A.E. Faraggi, C. Kounnas and J. Rizos, *Nucl. Phys.* **B774** (2007) 208;
C. Angelantonj, A.E. Faraggi and M. Tsulaia, JHEP **1**007, (2010) 004.

[20] T, Catelin–Jullien, A.E. Faraggi, C. Kounnas and J. Rizos, *Nucl. Phys.* **B812** (2009) 103;
A.E. Faraggi, I. Florakis, T. Mohaupt and M. Tsulaia, *Nucl. Phys.* **B848** (2011) 332.

[21] G.B. Cleaver and A.E. Faraggi, *Int. J. Mod. Phys.* **A14** (1999) 2335.

[22] M. Dine, N. Seiberg and E. Witten, *Nucl. Phys.* **B289** (1987) 589;
J.J. Atick, L.J. Dixon and A. Sen, *Phys. Rev.* **D292** (1987) 109.

[23] J. Pati and A. Salam, *Phys. Rev.* **D10** (1974) 275.

[24] A.E. Faraggi, *Phys. Lett.* **B245** (1990) 435;
A.E. Faraggi and E. Halyo, *Phys. Lett.* **B307** (1993) 311;
C. Coriano and A.E. Faraggi, *Phys. Lett.* **B581** (2004) 99.

[25] A.E. Faraggi and M. Guzzi, *Eur. Phys. Jour.* **C75** (2015) 537.

[26] A.E. Faraggi, *Nucl. Phys.* **B428** (1994) 111; *Phys. Lett.* **B520** (2001) 337.

[27] A.E. Faraggi and V.M. Mehta, *Phys. Rev.* **D84** (2011) 086006; *Phys. Rev.* **D88** (2013) 025006;
P. Athanasopoulos, A.E. Faraggi and V. Mehta, *Phys. Rev.* **D89** (2014) 105023.

[28] E. Witten, *Nucl. Phys.* **B471** (1996) 135.

[29] C. Coriano, S. Chang and A.E. Faraggi, *Nucl. Phys.* **B477** (1996) 65;
C. Coriano, A.E. Faraggi and M. Plümacher, *Nucl. Phys.* **B614** (2001) 233.

[30] S.Kalara, J.L. Lopez and D.V. Nanopoulos, *Nucl. Phys.* **B353** (1991) 650;
J. Rizos and K. Tamvakis, *Phys. Lett.* **B262** (1991) 227;
A.E. Faraggi, *Nucl. Phys.* **B487** (1997) 55.

[31] For a partial list see *e.g.*:
K. Harigaya and Y. Nomura, arXiv:1512.04850;
A. Pilaftsis, arXiv:1512.04931;
R. Franceschini *et al.*, arXiv:1512.04933;
S. Di Chiara, L. Marzola and M. Raidal, arXiv:1512.04939;
S.D. McDermott, P. Meade and H. Ramani, arXiv:1512.05326;
J. Ellis *et al*, arXiv:1512.05327;
R.S. Gupta *et al*, arXiv:1512.05332;
Q.H. Cao *et al*, arXiv:1512.05542; arxiv:1512.08441;
A. Kobakhidze *et al*, arXiv:1512.05585;
J.M. No, V. Sanz and J. Setford, arXiv:1512.05700;
W. Chao, R. Huo and J.H. Yu, arXiv:1512.05738;
L. Bian, N. Chen, D. Liu and J. Shu, arXiv:1512.05759;
J. Chakrabortty *et al*, arXiv:1512.05767;
A. Falkowski, O. Slone and T. Volansky, arXiv:1512.05777;
D. Aloni *et al*, arXiv:1512.05778;
W. Chao, arXiv:1512.06297;
S. Chang, arXiv:1512.06426;
R. Ding, L. Huang, T. Li and B. Zhu, arXiv:1512.06560;

X.F. Han, L. Wang, arXiv:1512.06587;
T.F. Feng, X.Q. Li, H.B. Zhang and S.M. Zhao, arXiv:1512.06696;
F. Wang, L. Wu, J.M. Yang and M. Zhang, arXiv:1512.06715;
F.P. Huang, C.S. Li, Z.L. Liu and Y. Wang, arXiv:1512.06732;
M. Bauer and M. Neubert, arXiv:1512.06828;
M. Chala, M. Duerr, F. Kahlhoefer and K. Schmidt-Hoberg, arXiv:1512.06833;
S.M. Boucenna, S. Morisi and A. Vicente, arXiv:1512.06878;
C.W. Murphy, arXiv:1512.06976;
G.M. Pelaggi, A. Strumia and E. Vigiani, arXiv:1512.07225;
J. de Blas, J. Santiago and R. Vega-Morales, arXiv:1512.07229;
A. Belyaev *et al*, arXiv:1512.07242;
P.S.B. Dev and D. Teresi, arXiv:1512.07243;
K.M. Patel and P. Sharma, arXiv:1512.07468;
S. Chakraborty, A. Chakraborty and S. Raychaudhuri, arXiv:1512.07527;
W. Altmannshofer *et al*, arXiv:1512.07616;
B.C. Allanach, P.S.B. Dev, S.A. Renner and K. Sakurai, arXiv:1512.07645;
N. Craig, P. Draper, C. Kilic and S. Thomas, arXiv:1512.07733;
J.A. Casas, J.R. Espinosa and J.M. Moreno, arXiv:1512.07895;
L.J. Hall, K. Harigaya and Y. Nomura, arXiv:1512.07904;
A. Salvio and A. Mazumdar, arXiv:1512.08184;
F. Wang *et al*, arXiv:1512.08434;
X. J. Bi *et al.*, arXiv:1512.08497;
F. Goertz, J.F. Kamenik, A. Katz and M. Nardecchia, arXiv:1512.08500;
P.S.B. Dev, R.N. Mohapatra and Y. Zhang, arXiv:1512.08507;
S. Kanemura, N. Machida, S. Odori and T. Shindou, arXiv:1512.09053;
I. Low and J. Lykken, arXiv:1512.09089;
A.E.C. Hernndez, arXiv:1512.09092;
Y. Jiang, Y.Y. Li and T. Liu, arXiv:1512.09127;
K. Kaneta, S. Kang and H. S. Lee, arXiv:1512.09129;
L. Marzola *et al*, arXiv:1512.09136;
X.F. Han *et al*, arXiv:1601.00534;
W. Chao, arXiv:1601.00633;
T. Modak, S. Sadhukhan and R. Srivastava, arXiv:1601.00836;
F.F. Deppisch *et al*, arXiv:1601.00952;
R. Ding, Z.L. Han, Y. Liao and X. D. Ma, arXiv:1601.02714;
J. Shu and J. Yepes, arXiv:1601.06891.

[32] B. Dutta *et al*, arXiv:1512.05439;
B. Dutta *et al*, arXiv:1601.00866 [hep-ph];
A. Karozas, S.F. King, G.K. Leontaris and A.K. Meadowcroft, arXiv:1601.00640.

[33] A.E. Faraggi, *Phys. Lett.* **B398** (1997) 88.

[34] A.E. Faraggi, M. Paraskevas, J. Rizos and K. Tamvakis, *Phys. Rev.* **D90** (2014) 015036.