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750 GeV diphoton excess in a $U(1)$ hidden symmetry model

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Recent results from the experimental collaborations at LHC give hints of a resonance in the diphoton channel at an invariant mass of 750 GeV. We show that such a scalar resonance would be possible in an $U(1)$ extension of the standard model where the extended symmetry is hidden and yet to be discovered. We explore the possibilities of accommodating this excess by introducing a minimal extension to the matter content and highlight the parameter space that can accommodate the observed diphoton resonance in the model. The model also predicts new interesting signals that may be observed at the current LHC run.

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I. INTRODUCTION

Recent results from the ATLAS and CMS Collaborations have shown the data from LHC run II with center of mass energy $\sqrt{s} = 13$ TeV [1,2]. Interestingly, the ATLAS data shows an excess in the diphoton channel with 3.2 fb^{-1} data giving about 14 events (with selection efficiency 0.4 [3]) at an invariant mass of ~750 GeV. The local significance is slightly northward of 3.5σ. On a lesser significance of about 2.6σ, a similar feature is exhibited by the CMS data with integrated luminosity of 2.6 fb^{-1}, giving about 10 events, peaked at an invariant mass of 760 GeV. The above rates with aforementioned efficiency correspond to a rough order of magnitude cross section of ~10 fb for the $pp \to X \to \gamma\gamma$. Although this can be a mere fluctuation in the early observations of the data at the upgraded energy run of LHC, the fact that both the collaborations observe it makes it an intriguing prospect for new physics signals. This naturally has led to a plethora of ideas explaining the excess [3–86].

In this work we show that a simple extension to the standard model (SM) gauge symmetry with a minimal set of new particles can easily accommodate the excess without invoking a large enough scale for new physics. In addition the model predicts some interesting signals that could show up as more data is accumulated in the run II of LHC. We consider an extra hidden $U(1)$ symmetry [87] in which all the SM particles are neutral. Only new exotic quarks, and an electroweak (EW) singlet Higgs boson can couple to this extra $U(1)$ gauge boson and the $U(1)$ symmetry is broken at the EW scale by the vacuum expectation value (VEV) of the EW singlet Higgs boson. In addition to this we extend the spectrum further by introducing an extra scalar which is a singlet under SM as well as the extra $U(1)$ symmetry [88]. We show that this scalar can be easily used to accommodate the observed diphoton excess with all particles of the model having masses within the TeV scale. In addition, we highlight new exotic decay modes of the vectorlike quark in the model that could give interesting signals at the LHC as well as a light sub-TeV Z' not constrained by existing experimental constraints.

II. MODEL

The gauge symmetry in our model [87] is the usual SM: $SU(3)_C \times SU(2)_L \times U(1)_Y$ supplemented by an extra $U(1)$ symmetry, which we call $U(1)_X$. We introduce two exotic quarks xq_L and xq_R which are color triplets but singlets under the $SU(2)_L$ gauge symmetry. They carry charge under the $U(1)_Y$ which decides whether they mix with the up-type or down-type SM quarks. We denote the gauge boson for the $U(1)_X$ by Z'. We introduce a complex Higgs field S_1 which acquires a VEV v_1 and breaks the $U(1)_X$. Therefore this scalar is a color and EW singlet, and has a charge q' under the $U(1)_X$. We also introduce a real scalar S_2 which is a singlet under $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_X$.

The EW gauge interaction Lagrangian for the exotic xq quark is given by

$$\mathcal{L} = \overline{xq} i \not{\mathcal{D}} xq \tag{1}$$

where the covariant derivative is defined as

$$\mathcal{D}_\mu = \partial_\mu - i\frac{g'}{2} Y B_\mu - i g_X Y_X Z'_\mu, \tag{2}$$

and Y_X is the charge of the matter field under the new gauge group $U(1)_X$ while Z' represents the new gauge boson.

The scalar potential, with the usual SM Higgs doublet H, and two new scalars, namely, the EW singlet S_1 and the real singlet S_2, is given by

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$$V(H,S_1,S_2) = -\mu_H^2(H^\dagger H) - \mu_{S_1}^2(S_1^\dagger S_1) - \mu_{S_2}^2 S_2^2 + \lambda_H(H^\dagger H)^2 + \lambda_{HS_1}(H^\dagger H)(S_1^\dagger S_1) + \lambda_{S_1}(S_1^\dagger S_1)^2 + \lambda_{S_2}S_2^4 + \lambda_{HS_2}(H^\dagger H)S_2^2 + \lambda_{S_1S_2}(S_1^\dagger S_1)S_2^2 + \sigma_1 S_2^3 + \sigma_2(H^\dagger H)S_2 + \sigma_3(S_1^\dagger S_1)S_2 \quad (3)$$

where the parameters $\mu_H, \mu_{S_1}, \mu_{S_2}, \sigma_1, \sigma_2$, and σ_3 have mass dimensions while $\lambda_H, \lambda_{HS_1}, \lambda_{S_1}, \lambda_{S_2}, \lambda_{HS_2}$, and $\lambda_{S_1S_2}$ are real dimensionless couplings. The EW symmetry is spontaneously broken when the neutral component of the Higgs doublet H gets a VEV while the additional $U(1)$ symmetry gets broken through the VEV of S_1. Then, in the unitary gauge, we can write the H, S_1, and S_2 fields as

$$H = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v_h + \mathcal{H} \end{pmatrix}, \quad S_1 = \frac{1}{\sqrt{2}}(v_1 + \mathcal{S}_1), \quad S_2 = v_2 + \mathcal{S}_2 \quad (4)$$

where v_h, v_1, and v_2 are VEVs of corresponding scalar fields while $\mathcal{H}, \mathcal{S}_1$, and $\mathcal{S}_2$ are the physical scalars in the gauge basis. Note that the terms in the above scalar potential with coefficients ($\lambda_{HS_1}, \lambda_{HS_2}, \lambda_{S_1S_2}, \sigma_2$, and σ_3) lead to a mixing between the three physical neutral scalars in the gauge basis, which we then choose to call h, h_s, and s in the mass basis, once the fields have acquired the VEV. We discuss the minimization conditions on the scalar potential, including constraints on the various coupling parameters ($\mu_i, \lambda_i, \sigma_i$) and the corresponding mass matrix relevant for this work in the Appendix.

After the neutral scalar fields have acquired VEVs, the SM gauge bosons ($Z, W^\pm$) get mass through the symmetry breaking via $\langle H \rangle = v_h/\sqrt{2} \sim v_{\rm EW}$ and the Z' gets mass via $\langle S_1 \rangle = v_1/\sqrt{2}$. We can also write a mass term for the vectorlike quark,

$$\mathcal{L}_{\rm mass} = M_x \overline{xq}_L xq_R. \quad (5)$$

Note that the new exotic vectorlike quark xq has color, hypercharge, and an extra $U(1)_X$ interaction, but no $SU(2)_L$ interaction. Since this new xq quark is vectorlike with respect to both $U(1)_Y$ as well as $U(1)_X$, the model is anomaly free. Without any other interaction, the xq quark will be stable. As none of the SM particles are charged under the new $U(1)_X$ symmetry, the new symmetry remains hidden from the SM, provided the gauge-kinetic-mixing terms are strongly suppressed. However, its gauge quantum numbers allow flavor changing Yukawa interactions with the SM quarks via the singlet Higgs boson S_1.

$$\mathcal{L}_{\rm Y_{extra}} = Y_{xq}\overline{xq}_L q_{iR} S_1 + {\rm H.c.} \quad (6)$$

where q_{iR} can be either the up-type or down-type quarks depending on the hypercharge we assign to xq for the above Lagrangian to be hypercharge singlet. We consider only mixing with the third generation quarks such that the hypercharge of both xq_L and xq_R must be equal to that of either t_R or b_R. This also requires that the $U(1)_X$ charge (Y_X) for the exotic quark xq must satisfy $Y_X = q'$. Such a term in the Lagrangian leads to mixing between the top (bottom) quark with the new exotic vectorlike quark xq, giving rise to EW decay modes for the heavy quark. In addition we can also write interaction terms for the new scalar S_2 with the xq given by

$$\mathcal{L} = -f_X \overline{xq} xq S_2. \quad (7)$$

Note that the vectorlike quark gets a bare mass as well as a mass from its Yukawa interaction with the singlet Higgs S_2, once S_2 gets a VEV. Note that using the above Lagrangian, the mass eigenstates from the mixing matrix for the q and xq [where $q = t(b)$ and $xq = xt(xb)$] along with their left and right mixing angles (θ_L, θ_R) can be determined using biunitary transformations.

Expressing the gauge eigenstates for the mixing quarks as q^0 and xq^0, the mass matrix in the (q^0, xq^0) basis is given by

$$\mathcal{M} = \begin{pmatrix} y_q v_h/\sqrt{2} & 0 \\ Y_{xq} v_1/\sqrt{2} & M_{xq} \end{pmatrix}, \quad (8)$$

where y_q is the usual Yukawa coupling of the SM quark with the Higgs doublet H while $M_{xq} = M_x - f_X v_2$. This matrix can be diagonalized with a biunitary transformation $\mathcal{M}_{\rm diag} = \mathcal{O}_L \mathcal{M} \mathcal{O}_R^\dagger$, where $\mathcal{O}_L$ and $\mathcal{O}_R$ are unitary matrices which rotate the left-chiral and right-chiral gauge eigenstates to the mass eigenstates respectively. The interaction of the physical mass eigenstates (q, xq) can then be obtained by writing the gauge basis states as

$$q_i^0 = q_i \cos\theta_i + xq_i \sin\theta_i, \qquad xq_i^0 = -q_i \sin\theta_i + xq_i \cos\theta_i, \quad (9)$$

while the rotation matrices $\mathcal{O}_i$ are given by

$$\mathcal{R}_i = \begin{pmatrix} \cos\theta_i & \sin\theta_i \\ -\sin\theta_i & \cos\theta_i \end{pmatrix}, \quad \text{where } i = L, R. \quad (10)$$

The corresponding mixing angles for the left- and right-handed fields follow from diagonalizing the matrices $\mathcal{M}\mathcal{M}^\dagger$ and $\mathcal{M}^\dagger\mathcal{M}$.

For our purposes, we can safely assume the mixings to be very small. However, it must be noted that such mixing, although small, would still ensure that the vectorlike quarks decay to SM quarks and bosons, i.e., $xq \to q'W, qZ, qh$. As the mixing angles θ_L and θ_R are constrained[1] by

[1]A detailed description on vectorlike quarks and mixing can be found in Ref. [89].

observables involving t, b quarks, in interactions within the SM as well as the entries in the Cabibbo-Kobayashi-Maskawa matrix, the small values would help in avoiding any such constraints easily. We also note that the model has three neutral scalars which also mix when H, S_1, and S_2 acquire VEVs. We must make the 125 GeV Higgs [90,91] to be SM-like and therefore dominantly the doublet component which therefore is unaffected in its properties by the presence of the exotic quark and scalar singlets. We, however, can try and allow significant mixing amongst the singlet scalars (see the Appendix). For simplicity we shall restrict our choice of the parameter space in the model, such that the mixing angles remain small.

A few comments are in order here:

(i) A quick look at the scalar potential [Eq. (3)] tells us that the mixing between the singlets and the doublet is related once the minimization conditions are imposed, as discussed in the Appendix.

(ii) One must also note that when the real singlet S_2 does not get a VEV, it is not possible to make the singlet-doublet ($\mathcal{H}-\mathcal{S}_2$) mixing vanishingly small while making the two singlets ($\mathcal{S}_1-\mathcal{S}_2$) mix substantially, as the mixing terms in the off-diagonal entries in the mass-squared matrix ($\mathcal{M}^2_{13}, \mathcal{M}^2_{23}$) in Eq. (A3) are of equal strength [by Eq. (A1) as $v_2 = 0$ and $v_h \sim v_1$], written in the ($\mathcal{H}, \mathcal{S}_1, \mathcal{S}_2$) basis. Note that $\mathcal{H}-\mathcal{S}_1$ mixing is independent of this and can be made negligibly small.

(iii) One can therefore achieve almost minimal mixing of the doublet with singlets while large mixing between the two singlets, once S_2 gets a VEV.

III. ANALYSIS

Thus, in our framework, we consider the most simplified scenario where the $xq \equiv xt$ has the same hypercharge as t_R and therefore mixes with the top quark. Note that although the mixing angles (θ_L, θ_R) can be arbitrary and small, it ensures a mixing which shall make the xt decay. Again, small mixings, if allowed in the scalar sector between H, S_1 and S_2, also ensure that xt, which had a dominant coupling with S_2, now also couples to the different scalar mass eigenstates (h, s, and h_s). Here the h is identified to be the SM-like Higgs boson. Therefore the vectorlike quark (VLQ) can decay through several modes if kinematically allowed.

We must point out that the new $U(1)$ gauge boson mass is given by $M_{Z'} = g_X q' v_1$ where v_1 is the VEV of S_1 that breaks $U(1)_X$ and q' is the $U(1)_X$ quantum number of S_1. Since this Z' only couples to top quarks, it is not possible to produce this directly at colliders and therefore existing bounds on such a Z' are very weak. The possible production channels for such a top-phillic Z' would be via associated production with $t\bar{t}$, and $xt\bar{x}\bar{t}$ or it can have loop-induced productions:

$$pp \to t\bar{t}Z'; \qquad xt\overline{xt}Z'; \qquad Z' + j(\text{loop}). \tag{11}$$

Note that in the absence of any kinetic mixing of the new $U(1)$ with SM Z, the Z' will have a four-body decay

$$Z' \to bW^+\bar{b}W^-$$

when $2m_b + 2m_W < m_{Z'} < m_t + m_b + m_W$, while Z' will have the three-body decays

$$Z' \to t\bar{b}W^-, \qquad \bar{t}bW^+$$

when $m_t + m_b + m_W < m_{Z'} < 2m_t$. A detailed phenomenological account of such a top-phillic Z' can be found in Refs. [92,93]. Note that in our model, the Z' has an additional mode of production which may become significant for lighter VLQ masses as well as the strength of the $U(1)_X$ gauge coupling, g_X. So the Z' can be much lighter than the heavier scalar mass eigenstates s and h_s as well as the VLQ. Thus xt can have quite a few possible decay products through the channels:

$$xt \to bW^+, \quad th, \quad tZ, \quad ts, \quad th_s, \quad tZ' \tag{12}$$

provided the mass states of s, h_s, Z' are lighter than xt. The additional decay modes would lead to new signals for the VLQ which can be produced at the LHC through strong interactions. As the existing bounds on such VLQ rely on its decay via bW^+, th, tZ modes only [94,95], the additional decay modes are expected to dilute the existing bounds on their mass and therefore one can have significantly lighter toplike VLQ still allowed by the experimental data. Signals for a VLQ with new decay modes to light neutral scalar have been considered before, for example in Ref. [88]. However, as we want to scan over the VLQ mass to fit the diphoton excess, there would be regions of parameter space where the VLQ becomes lighter than some of the above mentioned states, namely, h_s, s, or Z' which would disallow its decay to them. Since we set the mass of s to be 750 GeV, lighter xq can still decay through the remaining channels listed in Eq. (12). A much detail analysis of the VLQ and Z' signal at LHC in this model, which is interesting in itself is planned as future work.

For our current analysis, we shall consider the spectrum where s is dominantly composed of S_2 with a mass of around 750 GeV. Just like the VLQ, the scalar s can also decay via new channels other than a pair of SM particles including the Higgs boson (h). Namely, the new modes can be summarized as

$$s \to Z'Z', \quad h_s h_s, \quad xt\bar{t}, \quad t\overline{xt}, \quad hh_s, \tag{13}$$

again the decay being possible, depending on the mass of the decay products. The important thing to note here is that s would decay to gluon pair as well as diphoton via one-loop diagrams very similar to the SM Higgs boson, with the dominant contribution coming from xt in the loop (since S_2 couples to xt directly with a coupling strength f_X, which can be large). Thus, the production of this 750 GeV scalar is determined by the mass of xt and the size of the coupling strength f_X. The other decay channels for s can help in increasing the decay width of the resonance.[2] The loop-induced decay of the s to the massless gauge boson pairs gg is given by the effective Lagrangian

$$\mathcal{L}_{sGG} = -\lambda_{sgg} s G_{\mu\nu} G^{\mu\nu} \tag{14}$$

where $\lambda_{sgg} = \alpha_s f_X F_{1/2}(\tau_{xq})/(16\pi M_{xq})$. We choose the definition of the loop function $F_{1/2}(\tau_{xq})$ as given in Ref. [96]. We neglect the mixing effects here which can be justified by assuming that they are small enough to be negligible for the production rates but relevant to ensure the new decay modes for s and xq. However, as the new decay modes can reduce the branching fractions of the $s \to \gamma\gamma$, in order to fit the excess data, the mixings would be constrained to a great extent. For example,

(i) As the $s \to Z'Z', hh_s, h_s h_s$ decays happen when $\mathcal{S}_1 - \mathcal{S}_2$ mix, this mixing has to be taken small when the above decays are kinematically allowed for lighter Z', h_s so as not to suppress the diphoton mode significantly. In the current analysis we shall assume this mixing to be suppressed. Note that $s \to hh$ is disallowed by our choice of negligible mixing of the singlet scalars with the doublet Higgs as discussed in the Appendix.

We use the above interaction to calculate the production rates for the scalar s at the LHC run II and show the dependence of the cross section on the mass M_{xq} normalized to the coupling f_X. To do this we have implemented the effective vertex given by Eq. (14) in the event generator CalcHEP [97] and also include running of the strong coupling constant α_s calculated at $m_s = 750$ GeV in our estimates. The rates for the process shown in Fig. 1 is then simply given by the product of the production rate multiplied to the diphoton branching fraction which is naively $\alpha_{em}^2/\alpha_s^2(m_s) \lesssim 0.7\%$ at most if no additional decay modes of s are present.

In Fig. 2 we plot the leading-order (LO) production cross section for s with mass $m_s = 750$ GeV through the gluon-gluon fusion at the LHC with $\sqrt{s} = 13$ TeV as a function of the VLQ mass (M_{xq}). The cross sections are shown to be normalized with the coupling strength squared given by f_X^2.

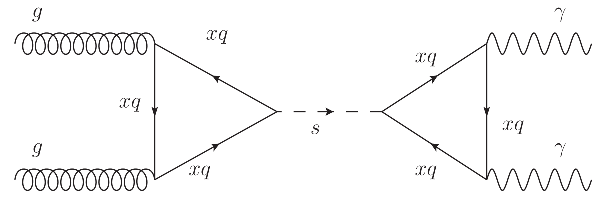

FIG. 1. Feynman diagram representing the diphoton resonant production through the scalar s at LHC.

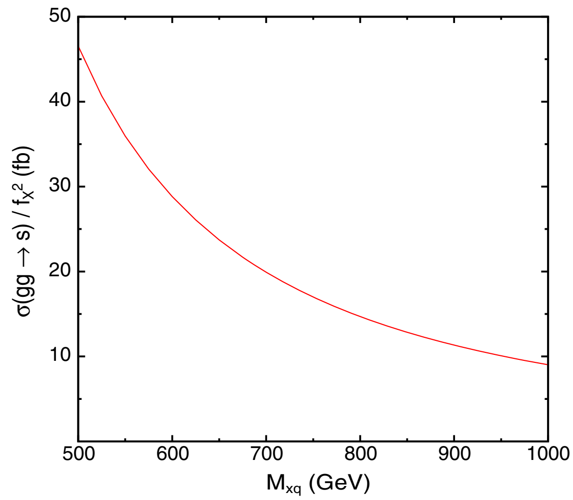

FIG. 2. The on-shell production cross section of s through gluon-gluon fusion at LHC with $\sqrt{s} = 13$ TeV as a function of the VLQ mass (M_{xq}). Note that we have normalized the cross section with the coupling strength squared (f_X^2).

We find that with only a single VLQ and without including any QCD corrections to the production, the production is as large as 46 fb for $M_{xq} = 500$ GeV and drops to about 10 fb when $M_{xq} = 1$ TeV with $f_X = 1$. We also find that with xt, the branching fraction for $s \to \gamma\gamma$ is about 0.6% which falls dramatically down to 0.04% if the VLQ is xb [due to the additional suppression from electric charge $(Q_d^2/Q_u^2)^2 \equiv 1/16$ to the partial width] for all values of M_{xq}. Note that the production cross section for the s is independent of this choice and therefore, quite clearly xt helps in enhancing the diphoton rates compared to xb. Assuming that $f_X \simeq \sqrt{4\pi}$ is taken at its perturbative limit, the production rates for s are enhanced by a factor of ∼12.57 which means that a resonant diphoton cross section with the toplike VLQ can be ≃10 fb with $M_{xt} \simeq 375$ GeV while achieving it with the bottom type VLQ will be clearly impossible. Of course one must note here that the QCD K factors for the $gg \to s$ production should not be very different from that of the SM Higgs. Including the QCD corrections can therefore simply double the production cross section ($K_{\rm NNLO} \simeq 2$), thus pushing the upper limit on M_{xq} to about 450 GeV to get ∼10 fb diphoton rate. For values of the VLQ mass less

[2]A wider resonance can also be realized (∼45 GeV) with both physical singlet masses (m_s, m_{h_s}) being close to around ∼750 GeV and separated by a small mass splitting. As both can contribute to the diphoton final state signal, it would be possible to explain the large width observed by the experimental collaborations without each scalar resonance being very wide itself.

than $m_s/2$, the tree-level decay mode, $s \to \bar{xq}xq$ opens up. This would completely dominate over all other decay channels making it practically impossible to fit the diphoton excess. Thus, $M_{xq} > m_s/2$ provides a lower bound to our choice of the VLQ mass. Quite clearly, one must resort to nonperturbative coupling strengths f_X for heavier VLQ mass to get the required cross section in the diphoton channel when including a single VLQ in the particle spectrum.

We, however, must point out that adding more generations of xq can easily alleviate this tension on the mass of the VLQ and the coupling f_X. Working within a single generation of VLQ, one can also include the bottom-like partner (xb) with the same $U(1)_X$ charge as the top partner (xt) as the minimal matter content in the model. This actually helps in increasing the production cross section by a factor of 4, assuming $M_{xb} = M_{xt}$ and $f_{xb} = f_{xt}$. However, the $s \to \gamma\gamma$ branching in this case drops to 0.25% which still effectively gives an enhancement of about 5/3 to the diphoton rates. This rate can be further enhanced by adding much lighter and less constrained vectorlike charged leptons ($x\tau$) that could enhance the photon branching significantly, thus easing the tension on the VLQ masses. In fact we find that for a single $x\tau$ with mass of about 400 GeV, the $s \to \gamma\gamma$ branching fraction peaks and goes up by a factor of ∼9 to about 4.5%, provided $f_{x\tau} \sim 3$, while $f_{xq} = 1$ and $M_{xq} = 600$ GeV. This would satisfy the 10 fb limit for the diphoton cross section with just xt as the VLQ with $M_{xt} \simeq 775(1050)$ GeV without (with) K factors, thus easily meeting the current limits on the VLQ mass. Notably, adding more vectorlike particles charged under the $U(1)_X$ gauge symmetry also enriches the Z' phenomenology of the model with additional production and decay channels. We leave these interesting possibilities to be taken up in a future work.

To show the relative dependence of including different set of vectorlike fermions in the particle spectrum, we plot the LO cross section of the diphoton signal at LHC with $\sqrt{s} = 13$ TeV in Fig. 3 as a function of the vectorlike fermion mass while keeping their Yukawa couplings with the singlet s to be equal ($f_X = 1$), which essentially represents the normalization condition ($1/f_X^2$) that we have put for the cross section. The M_{xq} has been varied between 400–1500 GeV. We can see that with just xb, it would be hard to achieve the observed diphoton signal. However, in all other scenarios, it is quite easy to observe a signal of 1–10 fb for the diphoton rate with perturbative values for f_X. As observed earlier, the inclusion of a vectorlike charged lepton with $M_{x\tau} = 400$ GeV allows the required signal rate, where VLQ masses are as high as 1.5 TeV. Note that a wide range of coupling and mass for the VLQ can easily accommodate the observed resonant signal. To summarize the plots, we note that the limit is around 700 GeV in the only toplike case. If both toplike and bottomlike VLQ are included, a little higher values in the

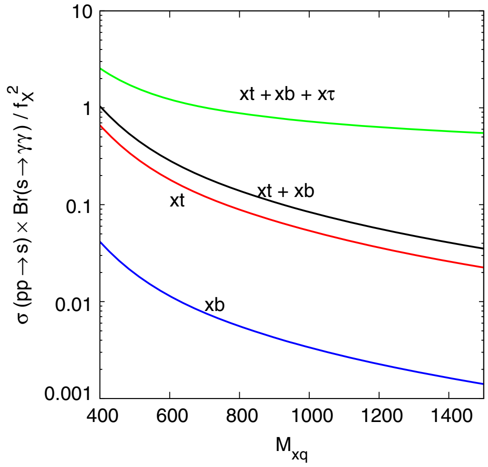

FIG. 3. The diphoton production rate for different exotic fermion scenarios with $\sqrt{s} = 13$ TeV as a function of the VLQ mass (M_{xq}). $M_{x\tau} = 400$ GeV has been taken.

mass of VLQ, i.e., around 900 GeV are achievable as the $gg \to s$ production cross section becomes 4 times that of only the toplike VLQ case. By including a 400 GeV vectorlike charged lepton, it is possible to push the vectorlike quark masses above 1 TeV. Independent of the exotic fermion content, f_X values below 0.5 are not allowed as long as the limit for the diphoton production rate is between 1–10 fb (at LO).

Thus, we find that within our model framework and a minimal extension of the matter particles by a single generation we can easily accommodate the diphoton excess without reverting to nonperturbative couplings or a very high new physics scale. However, as already mentioned earlier, in our model we have new decay modes [Eq. (12)] for the VLQ that not only relaxes the current limits on their masses but also leads to interesting signatures at the LHC which we leave for future analyses. We also expect that with more data collected by the experiment, the dijet resonance may show up at the same invariant mass for which the diphoton signal is observed (since the branching of $s \to gg$ can be significantly large for most of the parameter space), unless the other aforementioned decay modes of s become large. In addition, a very interesting signature in the model could be the production of light Z' through decays of the primarily produced VLQs.

IV. SUMMARY

In this work we show that a simple $U(1)$ extension to the SM gauge symmetry with a minimal set of new particles can easily accommodate the excess without invoking a large enough scale for new physics or nonperturbative couplings. We argue that with a high new physics scale, explaining the diphoton excess may lead to large nonperturbative coupling strengths for particle interactions. We show that a required low scale can be very easily realized within the context of our "hidden symmetry" model, thus keeping all couplings perturbative as well as complying with experimental constraints on the new physics scale.

We show that in our model the observed diphoton excess also highlights some new interesting signals that should show up as more data is accumulated in the run II of LHC. We perform a simplistic scan on the relevant parameters to show the compatibility of the resonant diphoton data with our model predictions. We also highlight the possibility of a very light Z' in the model with sub-TeV mass that can appear in decay cascades of the heavier particles such as the VLQ produced at the LHC. We leave the phenomenological analysis of such possibilities in our model as future work.

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APPENDIX: SCALAR MASS AND MIXINGS

We discuss the scalar potential of our model in some detail here. To find the minimum of the potential we use the following extremization conditions given by $\frac{\partial V}{\partial \mathcal{H}} = 0$, $\frac{\partial V}{\partial \mathcal{S}_1} = 0$, and $\frac{\partial V}{\partial \mathcal{S}_2} = 0$ which give us the following equations, respectively:

$$\begin{aligned} &\lambda_H v_h^3 + \frac{1}{2}\lambda_{HS_1} v_1^2 v_h + \lambda_{HS_2} v_2^2 v_h + \sigma_2 v_2 v_h - \mu_H^2 v_h = 0, \\ &\lambda_{S_1} v_1^3 + \frac{1}{2}\lambda_{HS_1} v_1 v_h^2 + \lambda_{S_1 S_2} v_1 v_2^2 + \sigma_3 v_1 v_2 - \mu_{S_1}^2 v_1 = 0, \\ &\lambda_{HS_2} v_2 v_h^2 + \lambda_{S_1 S_2} v_1^2 v_2 + 4\lambda_{S_2} v_2^3 + 3\sigma_1 v_2^2 + \frac{1}{2}(\sigma_2 v_h^2 + \sigma_3 v_1^2) - 2\mu_{S_2}^2 v_2 = 0. \end{aligned} \tag{A1}$$

Note that for the potential to be bounded from below we have

$$\lambda_H > 0, \qquad \lambda_{S_1} > 0, \qquad \lambda_{S_2} > 0. \tag{A2}$$

Using Eq. (A1) we can substitute for μ_H, μ_{S_1}, and μ_{S_2} in the scalar potential. Then the mass square matrix for the three neutral scalars in the gauge basis $(\mathcal{H}, \mathcal{S}_1, \mathcal{S}_2)$ becomes

$$\mathcal{M}^2 = \begin{pmatrix} 2\lambda_H v_h^2 & \lambda_{HS_1} v_1 v_h & (\sigma_2 + 2\lambda_{HS_2} v_2) v_h \\ \lambda_{HS_1} v_1 v_h & 2\lambda_{S_1} v_1^2 & (\sigma_3 + 2\lambda_{S_1 S_2} v_2) v_1 \\ (\sigma_2 + 2\lambda_{HS_2} v_2) v_h & (\sigma_3 + 2\lambda_{S_1 S_2} v_2) v_1 & \frac{1}{2v_2}(2(8\lambda_{S_2} v_2 + 3\sigma_1) v_2^2 - \sigma_2 v_h^2 - \sigma_3 v_1^2) \end{pmatrix}. \tag{A3}$$

For the point $(\mathcal{H} = 0, \mathcal{S}_1 = 0, \mathcal{S}_2 = 0)$ to be a minima of the potential, the matrix $\mathcal{M}^2$ should be positive definite, which is possible if its 3 upper left determinants are positive. The corresponding conditions are given below

$$2\lambda_H v_h^2 > 0; \qquad \begin{vmatrix} 2\lambda_H v_h^2 & \lambda_{HS_1} v_1 v_h \\ \lambda_{HS_1} v_1 v_h & 2\lambda_{S_1} v_1^2 \end{vmatrix} > 0 \Rightarrow 4\lambda_H \lambda_{S_1} - \lambda_{HS_1}^2 > 0; \quad |\mathcal{M}^2| > 0. \tag{A4}$$

For simplicity we have assumed that the mixing of $\mathcal{H}$ with $\mathcal{S}_1$ and $\mathcal{S}_2$ is vanishingly small and we shall set it to be zero. Note that such a choice not only imposes the condition that the scalar $\mathcal{H} \equiv h$ is a pure doublet component but also that it will have the exact properties of the SM Higgs boson with mass $m_h = \sqrt{2\lambda_H} v_h \simeq 125$ GeV. A quick look at the mass matrix then gives the conditions $\lambda_{HS_1} = 0$ and $\sigma_2 + 2\lambda_{HS_2} v_2 = 0$ for nonzero v_h and v_1.

We can now consider the two remaining singlet scalars independent of the doublet component $\mathcal{H}$. The reduced mass square matrix for the $\mathcal{S}_1$ and $\mathcal{S}_2$ sector becomes

$$M = \begin{pmatrix} 2\lambda_{S_1} v_1^2 & (\sigma_3 + 2\lambda_{S_1 S_2} v_2) v_1 \\ (\sigma_3 + 2\lambda_{S_1 S_2} v_2) v_1 & \frac{1}{2v_2}(2(8\lambda_{S_2} v_2 + 3\sigma_1) v_2^2 - \sigma_2 v_h^2 - \sigma_3 v_1^2) \end{pmatrix}. \tag{A5}$$

The fields $(\mathcal{S}_1, \mathcal{S}_2)$ can now be expressed in terms of mass eigenstates (h_s, s) where

$$\mathcal{S}_1 = h_s \cos\beta - s\sin\beta, \tag{A6}$$

$$\mathcal{S}_2 = h_s \sin\beta + s\cos\beta. \tag{A7}$$

The mixing angle β is given by

$$\tan 2\beta = \frac{2M_{12}}{M_{11} - M_{22}} \tag{A8}$$

and

$$\sin 2\beta = \frac{2M_{12}}{\sqrt{(M_{11} - M_{22})^2 + 4M_{12}^2}}, \tag{A9}$$

where M_{ij} is the $(i,j)^{th}$ element of M in Eq. (A5).

The mass eigenvalues for the two scalars s and h_s are then given by

$$m_1^2 = \frac{1}{2}(M_{11} + M_{22} - \sqrt{(M_{11} - M_{22})^2 + 4M_{12}^2}) \tag{A10}$$

and

$$m_2^2 = \frac{1}{2}(M_{11} + M_{22} + \sqrt{(M_{11} - M_{22})^2 + 4M_{12}^2}). \tag{A11}$$

For our analysis we have $m_h = 125$ GeV, $m_s = 750$ GeV while m_{h_s} is a free parameter which we can vary in the model. Note that it is possible to make h_s lighter than s as well as the vectorlike quarks by choosing parameters such that $M_{11} < M_{22}$. In the absence of mixing with the Higgs doublet, the h_s then decays to SM quarks through mixing of the VLQ with SM quarks.

Note that while the condition for the nonmixing of the doublet with either of the singlets may not forbid the couplings of the singlet s with h, it shall prevent the decay of s to any SM particle pair arising out of such mixings in the scalar sector. In fact the condition $\sigma_2 + 2\lambda_{HS_2} v_2 = 0$ leads to the exact cancellation of an interaction vertex between $h - h - s$ arising from the terms in the scalar potential given by $+\lambda_{HS_2}(H^\dagger H)S_2^2 + \sigma_2(H^\dagger H)S_2$. This is crucial in avoiding the possible decay of the 750 GeV singlet scalar to SM Higgs pair which is constrained by data [3]. Similarly, the decay of h_s to h pair is also forbidden due to the mixing suppression. The relevant vertices for the interactions within the scalar sector can be easily determined for the mass eigenstates and are given by

$$\begin{array}{llllll}
h & h_s & h_s & : & -2\lambda_{HS_2} v_h {s_\beta}^2 \\
h & h_s & s & : & -2\lambda_{HS_2} v_h c_\beta s_\beta \\
h & s & s & : & -2\lambda_{HS_2} v_h {c_\beta}^2 \\
h_s & h_s & s & : & (6{c_\beta}^2 s_\beta \lambda_{S_1} v_1 - 24 c_\beta {s_\beta}^2 \lambda_{S_2} v_2 - 2(2 - 3{s_\beta}^2) s_\beta \lambda_{S_1S_2} v_1 \\
 & & & : & -2(1 - 3{s_\beta}^2) c_\beta \lambda_{S_1S_2} v_2 - 6 c_\beta {s_\beta}^2 \sigma_1 - (1 - 3{s_\beta}^2) c_\beta \sigma_3) \\
h_s & s & s & : & -(6 c_\beta {s_\beta}^2 \lambda_{S_1} v_1 + 24 {c_\beta}^2 s_\beta \lambda_{S_2} v_2 + 2(1 - 3{s_\beta}^2) c_\beta \lambda_{S_1S_2} v_1 \\
 & & & : & -2(2 - 3{s_\beta}^2) s_\beta \lambda_{S_1S_2} v_2 + 6{c_\beta}^2 s_\beta \sigma_1 - (2 - 3{s_\beta}^2) s_\beta \sigma_3)
\end{array}$$

where $c_\beta = \cos\beta$ and $s_\beta = \sin\beta$.

A few benchmark points can be identified which give possibilities of a spectrum where $m_s \sim 750$ GeV while m_{h_s} is either heavier, lighter or has mass close to s. For example:

$$(\sigma_1, \sigma_3) = (-150, 65)\ \text{GeV}, \qquad (\lambda_{S_1}, \lambda_{S_2}, \lambda_{S_1S_2}, \lambda_{HS_2}) = (1, 0.2, -0.04, 0.05), \qquad (v_h, v_1, v_2) = (246, 750, 760)\ \text{GeV},$$

gives $m_{h_s} = 1.06$ TeV while $m_s = 749.1$ GeV with a very small mixing ($|\sin\beta| \sim 5.6 \times 10^{-3}$). Similarly,

$$(\sigma_1, \sigma_3) = (-150, 65)\ \text{GeV}, \qquad (\lambda_{S_1}, \lambda_{S_2}, \lambda_{S_1S_2}, \lambda_{HS_2}) = (1, 0.2, -0.05, 0.05), \qquad (v_h, v_1, v_2) = (246, 450, 750)\ \text{GeV},$$

gives $m_{h_s} = 636.4$ GeV while $m_s = 746.2$ GeV with again a suppressed mixing angle ($|\sin\beta| \sim 1.5 \times 10^{-2}$). However a slight variation in the model parameters also gives for

$$(\sigma_1, \sigma_3) = (-130, 90)\ \text{GeV}, \qquad (\lambda_{S_1}, \lambda_{S_2}, \lambda_{S_1S_2}, \lambda_{HS_2}) = (1, 0.19, -0.05, 0.1), \qquad (v_h, v_1, v_2) = (246, 531, 760)\ \text{GeV},$$

$m_{h_s} = 758.7$ GeV while $m_s = 747.8$ GeV with a not so suppressed mixing angle ($|\sin\beta| \sim 0.54$) which can give the possibility of two resonances that look like a single wide resonance, as observed by the ATLAS Collaboration.

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