Diphoton Signals from Colorless Hidden Quarkonia

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We show that quarkonia-like states of a hidden SU(N) gauge group can account for the 750 GeV diphoton excess observed by ATLAS and CMS, even with constituents carrying standard model hypercharge only. The required hypercharge is modest, varying between about 1.3–1.6 for strong SU(N) coupling, to 2–3 for weak SU(N) coupling, for N = 3, 4. This scenario predicts a variety of diphoton and multi-photon resonances, as well as photons from continuum pair production, and possibly exotic decays into standard model fermions, with no significant multi-jet resonances.

INTRODUCTION

If new particles are produced at the LHC, they have so far eluded detection, suggesting some suppression of their decays. In the presence of such suppression, bound states of these particles can be produced near threshold. Diphoton resonances from bound-state decay may then be the first harbingers of new physics. In this paper, we interpret the excess of diphoton events near 750 GeV reported by the ATLAS [1,2] and CMS [3,4] Collaborations as arising from the decay of a “quarkonium”-like state, η_{X}, bound by a hidden confining SU(N). Our model is minimal in that it assumes photoproduction as the main η_{X} production channel [7,9]. Thus, we take the η_{X} constituents X to carry hypercharge only, but no standard model (SM) SU(3)_c or SU(2)_L quantum numbers. This is in contrast to similar recent work which featured bound states with colored constituents [10–19]. While the production cross section is controlled by the hypercharge of the constituents, Y_{X}, and is proportional to Y_{X}^4, bound-state formation is controlled by the new strong force.

We will focus on the possibility that the constituents X are vector-like fermions in the (anti-)fundamental representation of the hidden SU(N) [4]. The lowest-lying J = 0 bound state is mainly produced via vector-boson fusion (VBF) processes, especially γγ fusion, while Drell–Yan (DY) production gives J = 1 bound states. Below we discuss two limiting cases, for which the bound-state properties can be readily estimated. In the first, the SU(N) is weakly coupled, such that the SU(N) confinement scale Λ_{h} is much smaller than the constituent mass m_{X}. The bound state is governed by the Coulomb SU(N) potential in this case, so that its properties and production cross section can be calculated perturbatively. In the second, the SU(N) is strongly coupled, and confinement effects are important. We will then rely on analogies with measured QCD quantities, mainly in the charmonium system, to infer the properties of η_{X}.

For simplicity, we restrict our attention to high values of Λ_{h}, such that η_{X} decays into hidden glueballs are kinematically forbidden. We stress that intermediate values of Λ_{h} can easily account for the diphoton signal, but we cannot give any quantitative estimate of branching ratios, binding energies etc. in this region.

In either case, with no light SU(N) flavors, the models obtained are quirk-like models [12,13,20,22] with uncolored quirks and relatively large Λ_{h}, which leads to microscopic strings. The distinguishing signatures of the models are a variety of multi-photon signatures, with dijet or multi-jet resonances absent or suppressed. In addition to diphoton events from bound-state decays, X–X pair production above threshold also gives multi-photon final states, either directly from annihilations into photons, or from annihilations into hidden glueballs which in turn decay into photons (or SM fermions).

We also consider models with additional light SU(N) flavors, which are SM-singlets, with couplings to X and SM fields. X–X pair production above threshold then gives a pair of X-hadrons which decay to SM fields through those couplings.

1 Scenarios with scalar constituents are possible as well, leading to different phenomenology. For example, the DY production of J = 1 state is suppressed due to the lack of spin.
This paper is organized as follows. We begin by reviewing the required cross section to diphotons and decay rates for the bound states. We describe two scenarios, roughly corresponding to small and large hidden couplings, in which diphoton decays of the lightest $XX$ bound state in the hidden sector can yield the correct signal. We discuss the phenomenology of the hidden sector, including the possibility of excited bound states and LHC signals from pair production of $X$–$\bar{X}$. Possible decays of $X$ itself are considered before we conclude.

**THE $\eta_X$ BOUND STATE AND DIPHOTON SIGNAL STRENGTH**

We begin with an overview of the different scales in the hidden sector. For concreteness, we assume that the 750 GeV resonance is a $1^1S_0$ bound state $\eta_X$ of a fundamental Dirac fermion $X$ of mass $m_X$ and hypercharge $Y_X$ (cf. Table I). The required cross section is \cite{16}

$$\sigma(pp \rightarrow \eta_X \rightarrow \gamma\gamma) \bigg|_{\sqrt{s}=13 \text{ TeV}} \approx 3-6 \text{ fb}.$$  

(1)

Since its constituents are not colored, the dominant production channel for $\eta_X$ is VBF, specifically photon fusion. Taking the result from Ref. 9, which includes the contributions from inelastic–inelastic, elastic–inelastic, and elastic–elastic processes, the total photoproduction signal strength at 13 TeV in the narrow width approximation is given by 3

$$\sigma_{p p \rightarrow \eta_X \rightarrow \gamma \gamma}^{13} = 5 \text{ fb} \left( \frac{\Gamma_{\text{tot}}}{21 \text{ MeV}} \right) \text{Br}(\eta_X \rightarrow \gamma\gamma)^2,$$  

(2)

where $\Gamma_{\text{tot}}$ is the total decay width of $\eta_X$. $Z\gamma$ and $ZZ$ production channels contribute an additional 8% to the cross section 23.

**Low $\Lambda_h$**

For small $\alpha_h$, the SU($N$) binding potential can be described in the Coulomb approximation 24

$$|R_{n0}(0)|^2 = C^2 \frac{\alpha_h^3 m_X^3}{2 n^3},$$  

(5)

where

$$C = \frac{1}{2} (C_1 + C_2 - C_\eta),$$  

(6)

and $C_\eta, C_1, C_2$ are the quadratic Casimirs of the bound state and its constituent particles. For constituents in the (anti-)fundamental representation and the singlet bound state of $\bar{X} X$, $C_\eta = 0$ and $C_1 = C_2 = \frac{N^2-1}{2N}$. Then

$$|R_{n0}(0)|^2 \sim \left( \frac{N^2-1}{2N} \right)^3 \frac{\alpha_h^3 M^3}{16 n^3},$$  

(7)

where we assumed that the bound-state energy $M \sim 2m_X$, and defined $\alpha_h$ as the hidden sector gauge coupling in the MS scheme, evaluated at the Bohr radius of the bound state

$$r_{\text{rms}} \approx a_0 = 2/(C \alpha_h m_X), \quad \alpha_h = \alpha_h(a_0^{-1})$$  

(8)

The decay rate into photons is then given by

$$\Gamma(\eta_X \rightarrow \gamma\gamma) = \frac{1}{4} N Y_X^2 \alpha^2 \left( \frac{N^2-1}{2N} \right)^3 \frac{\alpha_h^3}{M^3},$$  

(9)

where $\alpha = \alpha(M) \approx \alpha(M_Z) \approx 1/128$.  

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2 Quarkonia are often labelled in the form $n_{sLJ}$ in analogy with spectroscopic notation, where $S$, $L$, and $J$ are the spin, orbital, and total angular momentum quantum numbers. The radial excitation number $n_r$ is related to the principal quantum number $n$ by $n = n_r + L$. The radial excitation number and orbital angular momentum are used in parentheses for particle-like names of quarkonia states, e.g., $J/\psi(1S)$.

3 The cross-section at 8 TeV is about a factor of 2 smaller, and thus in tension with Run 1 diphoton searches. However, this ratio is subject to potentially large uncertainties 23 8.
The binding energy is given by (at leading order)

$$E_n = -\frac{1}{4\pi^2} \alpha h^2 n X,$$

so that the mass of this bound state is ($E_b = E_1$)

$$M = 2m_X + E_b = \left(2 - \frac{1}{4} \left( \frac{N^2 - 1}{2N} \right)^2 \alpha h^2 \right) m_X.$$

To calculate the signal strength, we also need the partial decay rates into the hidden sector, and into different SM particles. Since the resonance couples only to hypercharge, the decay rates into $Z\gamma$ and $ZZ$ are given by

$$\frac{\Gamma(\eta_X \rightarrow \{Z\gamma, ZZ\})}{\Gamma(\eta_X \rightarrow \gamma\gamma)} = \{2\tan^2 \theta_W, \tan^4 \theta_W\} \approx \{0.6, 0.08\},$$

where $\theta_W$ is the Weinberg angle with $\sin^2 \theta_W \approx 0.23$. Including these channels reduces the branching ratio to diphotons by about 40%.

Decays into hidden sector hadrons are dominated by annihilations into two hidden gluons, Eq. (4). With no light hidden flavors, these hadronize into hidden glueballs. For pure QCD, the lightest glueball has $J^{PC} = 0^{++}$ and mass $\sim 7\Lambda_{QCD}$ [25]. We assume the same scaling for the lightest glueball $G_h$, where the confinement scale $\Lambda_h$ is given at one-loop order by

$$\Lambda_h \sim m_X \exp \left( \frac{-2\pi}{b_0 \alpha h(m_X)} \right),$$

with $b_0 = \frac{11}{3} N - \frac{2}{3} N_F$, where $N_F$ is the number of light fermion flavors. $G_h$ mainly decays to photons through loops of $X$, with lifetimes estimated for $N = 3$ (see, e.g., Refs. 26–28),

$$\Gamma(0^{++}_h \rightarrow \gamma\gamma) \approx \frac{Y_A^4 \alpha^2}{64\pi^3} \frac{M_{G_h}^3}{m_X^6} \left( \frac{3M_{G_h}^2}{60m_X^3} \right)^2.$$

For fixed $N$, we are therefore left with $Y_X$ and $\alpha h$ as free parameters. In Figs. 2 and 3, we show the required diphoton signal strength in the Coulomb regime as a blue band in the $Y_X–\alpha h$ plane for $N = 3$ and $N = 4$, respectively, with $N_F = 0$. The plots are truncated at $\alpha h = 0.25$, where the Coulomb approximation can no longer be trusted (see discussion below). Also indicated on the $N = 3$ plot is the range of $Y_X$ preferred for the diphoton signal strength in the “High-$\alpha h$” scenario, which is discussed in the next subsection. As expected, lower values of $Y_X$ are preferred in this case compared to the Coulomb regime, since the enhancement factor $|R(0)|^2$ grows with the hidden coupling. In between the two regions, phenomenological potentials, which interpolate between the Coulomb and confining regimes, can be used to describe the bound states (see, e.g., Ref. 29). Clearly, there are large uncertainties in these calculations. Thus, for example, for stoponium, a recent lattice calculation finds that potential models [30] underestimate $|R(0)|^2$ by a factor of about 4 [31]. In any case,
the effect of these modifications would be to increase the production cross section of $\eta_X$, so that the blue region would shift to smaller values of $Y_X$ above a given $\alpha_h$.

The vertical red (dashed) lines denote the lightest glueball mass in GeV; here, we determine $\Lambda_h$ by solving the two-loop renormalization group (RG) equations for the boundary condition $\alpha_h(\Lambda_h) = 4\pi$. The orange (dotted) lines give the ratio $\Gamma(\eta_X\to\gamma\gamma)/\Gamma(\eta_X\to g_hg_h)$ from Eq. (14). The grey-filled regions in Figs. 2 and 3 indicate the lifetime of the lightest glueball according to Eq. (14), ranging from $>1$ s in the dark grey region to prompt decays in the white region. The dark grey region is excluded by Big Bang nucleosynthesis (BBN) since energetic photons from glueball decays would dissociate nuclei. The binding energy is less than $O(10\text{ GeV})$ in this range: we have included the contribution of the hypercharge Coulomb potential to the binding energy.$^4$ The solid black lines in the figures indicate values of $Y_X$, $\bar{\sigma}_h$ for which the contributions to the binding energy from hypercharge and the hidden SU($N$) are equal. For $Y_X$ between 2–2.5, we obtain the diphoton signal for $\bar{\sigma}_h > 0.22$ (0.18) for $N = 3$ (4), which corresponds to $\alpha_h(M) > 0.1$ (0.08).

For very small values of $\alpha_h$, the cross section becomes independent of $\bar{\sigma}_h$ and the hypercharge Coulomb potential is dominant in creating the bound state. The ratio of enhancement factors and binding energies are,

$$\frac{|R(0)|^2_{em}}{|R(0)|^2_h} \sim \left( \frac{E_{h,em}}{E_{h,h}} \right)^{3/2} \sim \left( \frac{2Y_X^2\alpha}{N\alpha_h} \right)^3.$$  \hspace{1cm} (15)

Above the black lines in Figs. 2 and 3, the hypercharge Coulomb interaction is dominant. In particular, we can use our results to estimate whether a purely hypercharged $X$ can account for the signal. This requires $Y_X \sim 4$, in contrast with larger values found in Ref. [22] (taking into account the different multiplicity $N$ in our model). Finally, we note that for $Y_X = 4$ and $N = 3$ ($N = 4$), the hypercharge coupling $g_Y$ becomes non-perturbative at around 2000 (300) TeV due to additional running from $X$.

For large $\bar{\sigma}_h$, we can understand the flattening of the blue signal region in Figs. 2 and 3 from Eq. (16) as follows. It is instructive to consider how the total rate into photons scales with the SU($N$) parameters and $Y_X$. From Eqs. (2), (3), and (5), and neglecting phase space factors due to the hidden glueball mass, we have

$$\sigma(pp\to\eta_X\to\gamma\gamma) \propto Y_X^8 \left( \frac{\bar{\sigma}_h}{\alpha_h(M)} \right)^2 \bar{\sigma}_h,$$  \hspace{1cm} (16)

which grows very fast with $Y_X$, and approximately linearly with the hidden coupling. In Eq. (8), we see that $a_0^{-1}$ grows linearly with $\bar{\sigma}_h$; therefore, for large $\bar{\sigma}_h$, the hierarchy between $a_0^{-1}$ and $M$ itself is small, so $\bar{\sigma}_h/\alpha_h(M)$ is approximately constant. Hence, to maintain a fixed cross section, $Y_X$ needs to change only by a small amount to compensate for a given change in $\bar{\sigma}_h$.

As $\alpha_h$ increases, confinement effects become important, and the Coulomb approximation becomes inadequate. Roughly speaking, this approximation is valid when the Bohr radius of the bound state is larger than the confinement scale$^4$

$$\Lambda_h \ll a_0^{-1} = \frac{C}{2} \bar{\sigma}_h m_X = \left( \frac{N^2 - 1}{4N} \right) \bar{\sigma}_h m_X. \hspace{1cm} (17)$$

Furthermore, the radial wavefunction at the origin $|R(0)|^2$ and the binding energy as given in Eqs. (5) and (10) are the leading-order results. Higher-order corrections to the binding energy were calculated in [33]. To obtain another estimate of the validity of the perturbative expansion, we take their result for the next-to-leading order (NLO) correction for zero light flavors,

$$E_h = E_1 \left[ 1 + \frac{\bar{\sigma}_h}{\pi} N \times 2.85 + O\left( \frac{\bar{\sigma}_h/\pi}{2} \right) \right], \hspace{1cm} (18)$$

which implies that we need $\bar{\sigma}_h < 0.25$ in order to trust the perturbative expansion. We therefore truncate the plots at $\bar{\sigma}_h = 0.25$. For both Figs. 2 and 3, Eq. (17) is satisfied in this region.

The above discussion assumed no light SU($N$) flavors for concreteness, but it can be simply extended to $N_F > 0$ light flavors with mass of order $\Lambda_h$. In this case, $\eta_X$ can decay to hidden hadrons as well, whose mass is probably of order $\Lambda_h$, and lighter than the glueballs. Furthermore, $\Lambda_h$ is smaller for a given $\alpha_h$ because of the slower running. For a single flavor, this is a mild effect, and the results of Figs. 2 and 3 will remain unchanged. In fact, the region in which the Coulomb approximation is valid will be wider in this case, allowing for larger $\alpha_h$.

**High $\Lambda_h$**

We now turn to consider larger SU($N$) couplings, for which the Coulomb approximation fails. In this scenario, $|R(0)|^2$ (and hence the enhancement of bound-state production) and the binding energy are sensitive to confinement effects, and are therefore large. In addition to the diphoton rates, $\eta_X$ will have large decay rates into hidden hadrons. Whether or not these are consistent with

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$^4$ Including the hypercharge contribution in the binding energy and the wavefunction at the origin involves the substitution $(C\bar{\sigma}_h)^n \to (C\bar{\sigma}_h + Y_X^2\alpha)^n$ in Eqs. (8) and (16).

$^5$ The potential approximation is valid for $\eta_h(1S)$, the lowest state of bottomonium, which has a separation scale $a_0 \sim 0.2$ fm compared to $\Lambda_{QCD}^2 \sim 1$ fm [29].
In the charmonium case, the mass of the $1^1S_0$ bound state $\eta_c(1S)$ is $M_{\eta_c} \sim 3$ GeV. The major difference in QCD is the existence of $N_F = 3$ light quarks below $\Lambda_{QCD}$. In this case, the $D$ mesons are the effective constituents of charmonium with $M_D \sim 1.85$ GeV for the lightest $D$ meson; therefore, the binding energy is approximately $0.6 - 0.7$ GeV, or about 20% of the mass of $\eta_c(1S)$. We then expect $m_X \sim 400 - 450$ GeV.

To obtain the decay width \( \Gamma(\eta_c \rightarrow \gamma\gamma) \), we can perform a similar scaling with charmonium. Using $\Gamma_{\text{tot}}/M$ for $\eta_c(1S)$ listed in Table II and $\text{Br}(\eta_c \rightarrow \gamma\gamma) = (1.59 \pm 0.12) \times 10^{-4}$, both taken from Ref. [34], we obtain

\[
\frac{\Gamma(\eta_c \rightarrow \gamma\gamma)}{M} = \frac{\Gamma(\eta_c \rightarrow \gamma\gamma)}{M_{\eta_c}} \left( \frac{Y_X}{Q_c} \right)^4
\]

\[
= 1.75 \times 10^{-6} \left( \frac{Y_X}{Q_c} \right)^4,
\]

where $Q_c = 2/3$ is the electromagnetic charge of the charm quark. Substituting the above for $\Gamma_{\text{tot}}$ and with $\text{Br} \sim 0.6$ in Eq. (2), we thus see that the diphoton excess can be accounted for $1.3 \lesssim Y_X \lesssim 1.6$, and for a wide range of $\alpha_s$ (and equivalently, of $\Lambda_h$).

The range is shown in the hashed region on the right-hand side of Fig. 2. Some representative values of these scales are sketched in Fig. 4, with the $X$ mass between 400–450 GeV, $\Lambda_h$ is between 100–150 GeV, and the glueballs between 600 GeV and 1 TeV, reflecting the large uncertainties in the glueball mass.

The large binding energy in this case allows for the possibility of additional bound states contributing to the signal. We will elaborate on this possibility below.

### Table III. Summary of branching ratios to radiative decays for $J/\psi(1S)$ and $\Upsilon(1S)$, from Ref. [34]. Decays in the last column proceed through a virtual photon.

| $J/\psi(1S)$ | $\text{Br}(1^1S_1 \rightarrow (gg\gamma\gamma))$ | $\text{Br}(1^1S_1 \rightarrow (had, \ell^+\ell^-))$ |
|--------------|---------------------------------|---------------------------------|
| $0.641, 0.088$ | $(0.135, 0.119)$ |
| $0.817, 0.022$ | $(\sim 0, 0.075)$ |

### Phenomenology: Other LHC Signals

#### Additional bound states

So far we considered threshold production of the $1^1S_0$ state $\eta_X$. As mentioned above, additional bound states of different masses may contribute to the diphoton signal. For example, in the charmonium system, $J/\psi(1S)$ decays to $\gamma + \eta_c(1S)$ with a branching ratio $0.017 \pm 0.004$. The hyperfine mass splitting between the $\eta_c(1S)$ and the $J/\psi(1S)$ is approximately $113$ MeV, or about $3.8\%$ of the $\eta_c(1S)$ mass. Scaling this splitting to $M$ yields a mass for the $1^1S_1$ bound state, which we call $T_X(1S)$ in analogy

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\[\text{Note that we cannot add light SU(N) flavors in this case, since the resulting mesons would provide new, dangerous decay channels for } \eta_X.\]
to bottomonium, that is 30 GeV above $M_f$. Therefore, the process $\Upsilon_X (1S) \rightarrow \eta_X \gamma$, which gives a relatively soft $\gamma$, can contribute to the diphoton signal. This process also contributes to the $\eta_X$ diphoton resonance. In the Low-$\Lambda_h$ scenarios, the hyperfine splitting of the 1S state in the Coulomb approximation is given by

$$\Delta E_{HF} \equiv E_b(1^3S_1) - E_b(1^1S_0) = \frac{1}{3} C^4 \tau^4 \rho m_X, \quad (21)$$

and can be $O(10)$ GeV.

At threshold, $\eta_X$ is produced through VBF while $\Upsilon_X$ is produced through DY. For $N_F = 0$, the possible $\eta_X$ and $\Upsilon_X$ decay channels are

$$\begin{align*}
\eta_X &\rightarrow VV, VG, G_h G_h, \\
\Upsilon_X &\rightarrow f \bar{f}, VVV, \eta_X V, G_h V, G_h G_h, \\
\end{align*} \quad (22)$$

where $V$ is a photon or a $Z$-boson, $f$ denotes a charged SM fermion, and $G_h$ is the excited 1$^+$ hidden glueball.

Among these, the process $\Upsilon_X \rightarrow l^+ l^-$ is severely constrained by LHC dilepton resonance searches $^{[36-39]}$, where $l = e, \mu$. These constraints imply

$$\sigma(pp \rightarrow \eta_X \rightarrow \gamma \gamma) < \frac{1.3}{K_{\eta\eta}} \text{fb} \cdot Y_X^2 \left( \frac{C_{\gamma\gamma}}{78} \right) \frac{\text{Br}(\eta_X \rightarrow \gamma \gamma)}{\text{Br}(\Upsilon_X \rightarrow f \bar{f})}, \quad (23)$$

In the above, we have employed the bound $\sigma(pp \rightarrow \Upsilon_X \rightarrow l^+ l^-) < 1.2 \text{fb}$ at the 8 TeV LHC $^{[38]}$. The relevant parton luminosities are taken from Ref. $^{[40]}$, except for $C_{\gamma\gamma}$, which is extracted from Eq. (2), and we set $K_{\eta\eta} = 1$. In the Low-$\Lambda_h$ region, where we can calculate the different $\Upsilon_X$ decay rates, we find that, for $N_F = 0$ and setting $K_{\eta\eta} = 1$, the resulting diphoton signal is smaller than 3 fb for $\tau_h \gtrsim 0.2$. For larger values of $\tau_h$, higher values of $\pi_h$ would be allowed, since $\Gamma(\Upsilon_X \rightarrow 3g_h)$ would increase as $\alpha_h(m_X)$ increases. In the High-$\Lambda_h$ case, the decay $\Upsilon_X \rightarrow G_h \gamma$ is open (see Fig. 4); however, to satisfy the constraint from $\Upsilon_X$ production, $\Gamma(\Upsilon_X \rightarrow G_h \gamma)$ must be enhanced by a factor of around 10 compared to the perturbative $\Gamma(\Upsilon_X \rightarrow g_h g_h \gamma)$.

Finally, radial excitations of $\eta_X$ may be produced as well. If these decay directly to diphotons (rather than to $\eta_X$ plus soft photons), they could lead to a broad diphoton peak, of width $\lesssim E_b$ $^{[18]}$. In order to obtain the $\sim 45$ GeV width favored by some analyses $^{[1]}$, we would require intermediate values for the couplings in between the Low-$\Lambda_h$ and High-$\Lambda_h$ regions.

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$^7$ This is consistent with the result expected from the confining potential, $\sim \lambda^2/M$.

$^8$ Production of more excited states (e.g., with $L > 0, n > 1$) is suppressed by their wavefunctions.

$^9$ This constraint assumes $\sigma(pp \rightarrow \Upsilon_X \rightarrow l^+ l^-) \equiv \sigma(pp \rightarrow \Upsilon_X \rightarrow e^+ e^-) = \sigma(pp \rightarrow \Upsilon_X \rightarrow \mu^+ \mu^-)$.

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**X–X pair production above threshold**

Here the signatures crucially depend on the presence of light SU($N$) flavors. For $N_F = 0$, the models are essentially quirk models, with the quirks carrying hypercharge only. The $X–X$ pairs will form a string of length $^{[20]}

$$r \sim \frac{E}{A^2}, \quad (24)$$

where $E$ is the quirk-pair energy. In all our models, this string is microscopic:

$$r \ll 1 \mu m. \quad (25)$$

The quirk pair will then promptly annihilate into photons. Note that the $X–X$ pair is produced with $L > 0$, and in principle, it could first relax to the bound state $\eta_X$, which would subsequently decay to diphotons, contributing to the diphoton resonance from $\eta_X$ threshold-production. To lose angular momentum, the excited state must radiate photons or glueballs. The latter process is kinematically suppressed in the range of relatively large $\Lambda_h$ we consider, so the only option is photon radiation. If this energy loss process is efficient, the signal would be relatively soft photons from the relaxation process, plus a diphoton resonance at 750 GeV. This may greatly enhance the resonance, since DY production is larger than photon fusion. However, with tight photon isolation cuts, these processes may be vetoed because of the presence of additional photons. The other possibility is that the radiation loss is not efficient, in which case $X$ and $\bar{X}$ annihilate without losing significant energy. This seems more plausible in our models, given the tight binding of the quirk pair by the microscopic string. The $X–\bar{X}$ pair can then go into either glueballs or photons, and the resulting photon invariant mass is simply the $X–\bar{X}$ invariant mass in this case.

Finally, the glueballs annihilate into SM particles, mainly photons, through $X$ loops as mentioned above. These events would then have four or more photons, with di- (or tri-) photon peaks at the glueball masses.

We now consider models with additional SU($N$) light flavors, with mass of order $\Lambda_h$. As explained above, this is only viable in the Low-$\Lambda_h$ case. The glueball lifetime is hardly affected, while the collider phenomenology as well as cosmology are different since string breaking is no longer suppressed. Consider an additional SM-singlet, SU($N$) fundamental field $S$, which can be either a scalar or a vector-like fermion. With no new couplings to SM fields, continuum $X–\bar{X}$ pair production would be followed by hadronization into $X–\bar{S}$ mesons. We refer to these mesons as $\xi_S$. The lightest $\xi_S$ is charged and stable. New couplings involving $X$, $S$, and SM fields must be introduced to mediate $\xi_S$ decays. This restricts $Y_X$ to integer values. The possible couplings are summarized in Table IV.
TABLE IV. Lowest-order operators mediating $\xi_S$ decays. $\phi_S$ ($\chi_S$) denotes a scalar (vector-like fermion) which is a SM singlet and SU(N) (anti-)fundamental. SM flavor indices are omitted for simplicity. If $Y_X < 0$, replace hidden sector fields by their conjugates.

| $Y_X$ | scalar $\phi_S$ | vector-like fermion $\chi_S$ |
|-------|----------------|-----------------------------|
| 1     | $\phi_S(\bar{X}e^c)$ | $\bar{X}e e^c$ |
| 2     | $\phi_S^0 X u^c u^e \bar{e}^c$, $\phi_S(\bar{X}e^c)(I)^*$ | $\bar{X}e e^c$ |
| 3     | $\phi_S \bar{X} e^c \bar{e}^c e^c$ | none up to dim-7 |

These decays provide various exotic signatures, with $\xi_S$ pair production followed by $\xi_S$ decay to $2\ell$, $3\ell$, or 3 jets, with $\ell = e, \mu$ or $\tau$. Events with $2\ell + E_T$ are possible too. With the exception of the scalar coupling for $Y_X = 1$, these couplings are non-renormalizable, and therefore naturally small. Thus, bound-state formation is still important, and $X$ particles indeed hadronize before decaying.

Still, these operators can give sufficiently high rates to evade the stringent constraints on long-lived (on detector scales) charged particles [41–43] if the scale by which they are suppressed is 10 TeV or higher, depending on the operator. Thus for example, for $Y_X = 2$, even the operator $(\lambda/M_{\text{high}}^2)\phi_S(\bar{X}e^c)(I)^*$ gives

$$c\tau \sim 30 \mu m \cdot \frac{1}{\lambda^2} \left( \frac{M_{\text{high}}}{10 \text{ TeV}} \right)^6 \left( \frac{1 \text{ GeV}}{\Lambda_h} \right)^2 \left( \frac{375 \text{ GeV}}{M_{\xi_S}} \right)^5,$$

while the analogous operator with $\chi_S$, allows for a higher $M_{\text{high}}$ since it is dimension-4 only.

CONCLUSIONS

In this paper, we considered the possibility that the observed diphoton excess is due to a quarkonium-like bound state, $\eta_X$, of a hidden SU(N), with fermionic constituents carrying SM hypercharge only. The production and decay of this bound state are controlled by two parameters: $Y^{\Delta \alpha}_X$, which sets the coupling strength to photons, and $C_{\alpha h}$, which controls the coupling to hidden gluons. These scenarios lead to a variety of multi-photon signals, and possibly exotic decays to SM fermions. Diphotons from photon-fusion production of $\eta_X$ are typically accompanied by forward jets, with hadronic activity in the central region suppressed.

In large parts of the parameter space, production of the $J = 1 \Upsilon_X$ bound state leads to dilepton or dijet resonances close to 750 GeV. Without additional hidden flavors or other dynamics, the bound Eq. [23] excludes the parameter space above $\sigma_{pp} \Upsilon_X \to l^+ l^-$ at the 13 TeV LHC is between 1–3 fb.

While we focused on a constituent fermion $X$ for concreteness, our results generalize trivially to the case of a scalar $X$. In particular, the blue curves of Figs. 2 and 3 are barely modified, since the scalar production cross sections are down by a factor of 2 compared to the fermion case, but the signal strength scales at least as the fourth power of $Y_X$. On the other hand, as mentioned in the Introduction, the production of $J = 1$ bound states is suppressed in this case.

More generally, the relation between $\eta_X$ and $\Upsilon_X$ production depends on the details of the model. For example, with $N_F$ flavors with masses somewhat below $m_X$, the renormalization group running of $\alpha_h$ between $m_X$ and the inverse Bohr radius is milder, resulting in larger $\sigma_{\eta_X}(m_X)$ for a given $\Lambda_h$. This enhances the rates of both $\Upsilon_X$ and $\eta_X$ to hidden gluons, with the former increasing more sharply. Second, in the presence of $S$–$X$ couplings to the SM, $X$ is unstable. If the decay width $\Gamma_X$ of $X$ is comparable to roughly half the bound state width, bound state production is depleted. Since $\Gamma_{\eta_X} \gg \Gamma_X$, the $\Upsilon_X$ production cross section is substantially reduced if $\Gamma_X \sim \Gamma_{\Upsilon_X} / 2$, whereas $\eta_X$ production is largely unaffected. Third, a hidden $Z'$ can mediate additional $\Upsilon_X$ decays to hidden flavors, while its effects on $\eta_X$ decays are milder.

Finally, $\eta_X$ production in association with additional photons or hidden gluons is quite generic in these models, whether its source is higher resonances or continuum quark-pair production for $N_F = 0$ models. Determining whether the diphoton resonance is accompanied by additional softer photons or missing energy is therefore crucial.

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