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Explaining the 750 GeV diphoton excess with a colored scalar charged under a new confining gauge interaction

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We consider a charged scalar particle χ of mass around 375 GeV charged under both $SU(3)_C$ and a new confining non-Abelian gauge interaction. After pair production, these interactions confine the exotic scalar into nonrelativistic bound states whose decays into photons can explain the 750 GeV diphoton excess observed at the LHC. Taking the new confining group to be SU(2), we find χ must carry an electric charge of $Q \sim [\frac{1}{2}, 1]$ to fit the data. Interestingly, we find that pair production of the scalars and the subsequent formation of the bound state dominates over direct bound state resonance production. This explanation is quite weakly constrained by current searches and data from the forthcoming run at the LHC will be able to probe our scenario more fully. In particular dijet, monojet, di-Higgs, and jet + photon searches may be the most promising discovery channels.

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I. INTRODUCTION

An excess of events containing two photons with invariant mass near 750 GeV has been observed in 13 TeV proton-proton collisions by the ATLAS and CMS collaborations [1,2]. The cross section $\sigma(pp \to \gamma\gamma)$ is estimated to be

$$\sigma(pp \to \gamma\gamma) = \begin{cases} (10 \pm 3)\ \text{fb} & \text{ATLAS} \\ (6 \pm 3)\ \text{fb} & \text{CMS} \end{cases} \tag{1}$$

and there is no evidence of any accompanying excess in the dilepton channel [3]. If we interpret this excess as the two photon decay of a single new particle of mass m then ATLAS data provide a hint of a large width: $\Gamma/m \sim 0.06$, while CMS data prefer a narrow width. Naturally, further data collected at the LHC should provide a clearer picture as to the nature of this excess.

There has been vast interest in the possibility that the diphoton excess results from physics beyond the standard model (SM). Most discussion has focused on models where the excess is due to a new scalar particle which subsequently decays into two photons; see e.g. [4] (for a recent discussion see also [5]). The possibility that the new scalar particle is a bound state of exotic charged fermions has also been considered, e.g. [6–10]. Here we consider the case that the 750 GeV state is a nonrelativistic bound state constituted by an exotic *scalar* particle χ and its antiparticle, charged under $SU(3)_C$ as well as a new unbroken non-Abelian gauge interaction. Having χ be a scalar rather than a fermion is not merely a matter of taste: In such a framework a fermionic χ would lead to the formation of bound states which (typically) decay to dileptons more often than to photons; a situation which is not favoured by the data.

The bound state, which we denote Π, can be produced through gluon-gluon fusion directly (i.e. at threshold $\sqrt{s_{gg}} \simeq M_\Pi$) or indirectly via $gg \to \chi^\dagger\chi \to \Pi +$ soft quanta (i.e. above Π threshold: $\sqrt{s_{gg}} > M_\Pi$). The indirect production mechanism can dominate the production of the bound state, which is an interesting feature of this kind of theory.

II. THE MODEL

We take the new confining unbroken gauge interaction to be $SU(N)$, and assume that, like $SU(3)_C$, it is asymptotically free and confining at low energies. However, the new $SU(N)$ dynamics is qualitatively different from QCD as all the matter particles [assumed to be in the fundamental representation of $SU(N)$] are taken to be much heavier than the confinement scale, Λ_N. In fact we here consider only one such matter particle, χ, so that $M_\chi \gg \Lambda_N$ is assumed. In this circumstance a $\chi^\dagger\chi$ pair produced at the LHC above the threshold $2M_\chi$ but below $4M_\chi$ cannot fragment into two jets. The $SU(N)$ string which connects them cannot break as there are no light $SU(N)$-charged states available. This is in contrast to heavy quark production in QCD where light quarks can be produced out of the vacuum enabling the color string to break. The produced $\chi^\dagger\chi$ pair can be viewed as a highly excited bound state, which deexcites by $SU(N)$-ball and soft glueball/pion emission [11].

With the new unbroken gauge interaction assumed to be $SU(N)$ the gauge symmetry of the SM is extended to

$$SU(3)_C \otimes SU(2)_L \otimes U(1)_Y \otimes SU(N). \tag{2}$$

This kind of theory can arise naturally in models which feature large color groups [12–14] and in models with

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leptonic color [15–17] but was also considered earlier by Okun [18]. The notation *quirks* for heavy particles charged under an unbroken gauge symmetry (where $M_\chi \gg \Lambda_N$) was introduced in [11] where the relevant phenomenology was examined in some detail in a particular model.[1] For convenience we borrow their nomenclature and call the new quantum number *hue* and the massless gauge bosons *huons* ($\mathcal{H}$).

The phenomenological signatures of the bound states (quirkonia) formed depend on whether the quirk is a fermion or boson. Here we assume that the quirk χ is a Lorentz scalar in light of previous work which indicated that bound states formed from a fermionic χ state would be expected to be observed at the LHC via decays of the spin 1 bound state into opposite-sign lepton pairs ($\ell^+\ell^-$) [11,17]. In fact, this appears to be a serious difficulty in attempts to interpret the 750 GeV state as a bound state of fermionic quirk particles (such as those of [7–9]). The detailed consideration of a scalar χ appears to have been largely overlooked,[2] perhaps due to the paucity of known elementary scalar particles. With the recent discovery of a Higgs-like scalar at 125 GeV [21,22] it is perhaps worth examining signatures of scalar quirk particles. In fact, we point out here that the two photon decay is the most important experimental signature of bound states formed from electrically charged scalar quirks. Furthermore this explanation is only weakly constrained by current data and thus appears to be a simple and plausible option for the new physics suggested by the observed diphoton excess.

III. EXPLAINING THE EXCESS

The scalar χ that we introduce transforms under the extended gauge group [Eq. (2)] as

$$\chi \sim (\mathbf{3}, \mathbf{1}, Y; \mathbf{N}), \tag{3}$$

where we use the normalization $Q = Y/2$. The possibility that χ also transforms nontrivially under $SU(2)_L$ is interesting; however, for the purposes of this paper we focus on the $SU(2)_L$ singlet case for definiteness. Since two-photon decays of nonrelativistic quirkonium will be assumed to be responsible for the diphoton excess observed at the LHC, the mass of χ will need to be around 375 GeV.

We have assumed that χ is charged under $SU(3)_C$ so that it can be produced at tree level through QCD-driven pair production. We present the production mechanisms in Fig. 1. To estimate the production cross section of the bound states, we first consider the indirect production mechanism which we expect to be dominant. Here, a $\chi^\dagger\chi$ pair is produced above threshold and deexcites emitting soft glueballs/pions and hueballs: $gg \to \chi^\dagger\chi \to \Pi +$ soft quanta. We first consider the case where the confinement scale of the new $SU(N)$ interaction is similar to that of QCD. What happens in this case can be adapted from the discussion in [11], where a fermionic quirk charged under an unbroken SU(2) gauge interaction was considered. As already briefly discussed in the Introduction, the $\chi^\dagger\chi$ pairs initially form a highly excited bound state, which subsequently deexcites in two stages. The first stage is the nonperturbative regime where the hue string is longer than Λ_N^{-1}. The second stage is characterized by a string scale significantly less than Λ_N^{-1}: the perturbative Coulomb region. Here the bound state can be characterized by the quantum numbers n and l. Deexcitation continues until quirkonium is in a lowly excited state with $l \leq 1$ and n. Imagine first that deexcitation continued until the ground state ($n = 1, l = 0$) is reached. Given we are considering χ to be a scalar, the quirkonium ground state, Π, will have spin 0, and is thus expected to decay into SM gauge bosons and huons. The cross section $\sigma(pp \to \Pi \to \gamma\gamma)$ in this case is then

$$\sigma(pp \to \gamma\gamma) \approx \sigma(pp \to \chi^\dagger\chi) \times \mathrm{Br}(\Pi \to \gamma\gamma). \tag{4}$$

Since production is governed by QCD interactions, we can use the values of the pair production cross sections for stops/sbottoms in the limit of decoupled squarks and gluinos [23]. For a χ mass of 375 GeV

$$\sigma(pp \to \chi^\dagger\chi) \approx \begin{cases} 2.6N \text{ pb} & \text{at 13 TeV} \\ 0.5N \text{ pb} & \text{at 8 TeV} \end{cases}. \tag{5}$$

The branching fraction is to leading order:

$$\mathrm{Br}(\Pi \to \gamma\gamma) \simeq \frac{3NQ^4\alpha^2}{\frac{2}{3}N\alpha_S^2 + \frac{3}{2}C_N\alpha_N^2 + 3NQ^4\alpha^2}, \tag{6}$$

where $C_N \equiv (N^2 - 1)/(2N)$, α_N is the new $SU(N)$ interaction strength and we have neglected the small contribution of $\Pi \to Z\gamma/ZZ$ to the total width. Equation (6) also neglects the decay to Higgs particles: $\Pi \to hh$, which arises from the Higgs potential portal term $\lambda_\chi \chi^\dagger\chi\phi^\dagger\phi$. Theoretically this rate is unconstrained given the dependence on the unknown parameter λ_χ, but could potentially be important. However, limits from resonant Higgs boson pair production derived from 13 TeV data, $\sigma(pp \to X \to hh \to bbbb) \lesssim 50$ fb at $M_X \approx 750$ GeV [24,25], imply that the Higgs decay channel must indeed be subdominant (cf. $\Pi \to gg, \mathcal{HH}$).

The renormalized gauge coupling constants in Eq. (6) are evaluated at the renormalization scale $\mu \sim M_\Pi/2$. Taking for instance the specific case of $N = 2$, $\alpha_N = \alpha_S \simeq 0.10$ (at $\mu \sim M_\Pi/2$) gives

$$\sigma(pp \to \gamma\gamma) \approx 5\left(\frac{Q}{1/2}\right)^4 \text{ fb} \quad 13 \text{ TeV}. \tag{7}$$

At $\sqrt{s} = 8$ TeV the cross section is around five times smaller. We present the cross section $\sigma(pp \to \Pi \to \gamma\gamma)$ for a range of masses M_Π and different combinations of Q and N in Fig. 2. The parameter choice $\alpha_N = \alpha_S$ and $\Lambda_N = \Lambda_{QCD}$

[1] Some other aspects of such models have been discussed over the years, including the possibility that the $SU(N)$ confining scale is low (∼keV), a situation which leads to macroscopic strings [19].

[2] The idea has been briefly mentioned in recent literature [9,20].

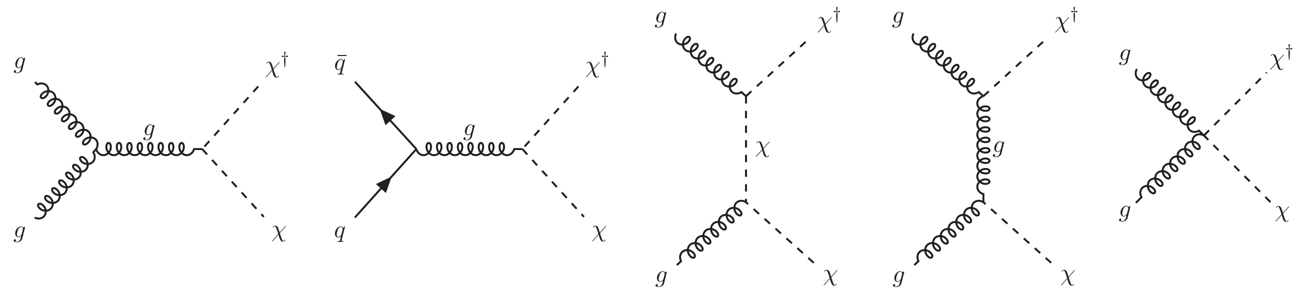

FIG. 1. Tree-level pair production mechanisms for the scalar quirk χ.

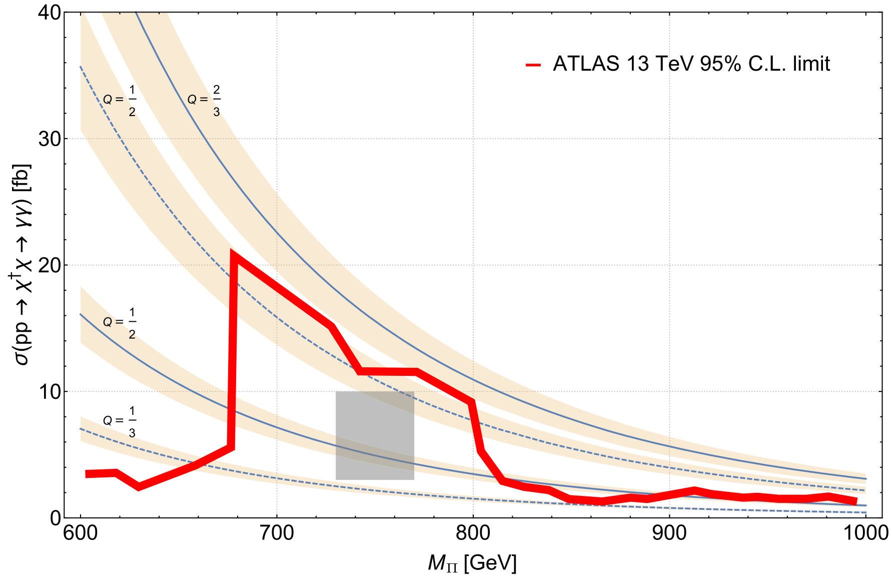

FIG. 2. The cross section $\sigma(pp \to \Pi \to \gamma\gamma)$ at 13 TeV for a range of quirkonium masses M_Π and charge assignments. Solid lines denote choices of $N = 2$ and dashed lines choices of $N = 5$. The rectangle represents the $\sigma \in [3, 10]$ fb indicative region accommodated by the ATLAS and CMS data. The solid red line is the ATLAS 13 TeV exclusion limit. Uncertainties reflect error associated with the parton distribution functions.

has been assumed. (The cross section is not highly sensitive to Λ_N, α_N so long as we are in the perturbative regime: $\Lambda_N \lesssim \Lambda_{\rm QCD}$.) Evidently, for $N = 2$, a χ with electric charge $Q \approx 1/2$ is produced at approximately the right rate to explain the diphoton excess.

In practice deexcitation of the produced quirkonium does not always continue until the ground state is reached. In this case annihilations of excited states can also contribute. However those with $l = 0$ will decay in the same way as the ground state. The only difference is that the excited states will have a slightly larger mass (which we will estimate in a moment) due to the change in the binding energy. This detail could be important as it can effectively enlarge the observed width. Annihilation of excited states with nonzero orbital angular momentum could in principle also be important; however, these are suppressed as the radial wavefunction vanishes at the origin: $R(0) = 0$ for $l \geq 1$. They are expected to deexcite predominately to $l = 0$ states rather than annihilate [11]. Nevertheless, for sufficiently large α_N the $l=1$ annihilations, $\Pi \to \mu^+\mu^-$ and $\Pi \to e^+e^-$, could potentially be observable.

The $l = 0$ excited states can be characterized by the quantum number n with binding energies:

$$\frac{E_n}{M_\Pi} = -\frac{1}{8n^2}\left[\frac{4}{3}\bar{\alpha}_S + C_N\bar{\alpha}_N + Q^2\bar{\alpha}\right]^2. \qquad (8)$$

The above formula was adapted from known results with quarkonium, e.g. [26] (and of course also the hydrogen atom). The coupling constants $\bar{\alpha}_S$, $\bar{\alpha}_N$, and $\bar{\alpha}$ are evaluated at a renomalization scale corresponding to the mean distance between the particles which is of order the Bohr radius: $a_0 = 4/[(4\bar{\alpha}_S/3 + C_N\bar{\alpha}_N + Q^2\bar{\alpha})M_\Pi]$. The bound state, described by the radial quantum number n has mass given by $M_\Pi(n) = 2M_\chi + E_n$. Considering as an example $N = 2$ and $\bar{\alpha}_N = \bar{\alpha}_S = 0.15$, $\bar{\alpha} = 1/137$ we find the mass difference between the $n = 1$ and $n = 2$ states to be $\Delta M = (E_1 - E_2) \approx 0.01 M_\Pi$. Larger mass splitings will be possible[3] if $\bar{\alpha}_N > \bar{\alpha}_S$, although it has been shown in the context of fermionic quirk models that the phenomenology is substantially altered in this regime [7]. In particular, the hueballs can become so heavy that the decays of the bound state into hueballs is kinematically forbidden.

In the above calculation of the bound state production cross section, we considered only the *indirect* production following pair production of $\chi^\dagger\chi$ above threshold. The bound state can also be produced directly: $gg \to \Pi$, where $\sqrt{s_{gg}} \approx M_\Pi$. The cross section of the ground state direct resonance production is

$$\sigma(pp \to \Pi)_{\rm DR} \approx \frac{C_{gg} K_{gg} \Gamma(\Pi \to gg)}{s M_\Pi}, \tag{9}$$

where C_{gg} is the appropriate parton luminosity coefficient and K_{gg} is the gluon NLO QCD K-factor. For $\sqrt{s} =$ 13 TeV we take $C_{gg} \approx 2137$ [4] and $K_{gg} = 1.6$ [27]. The partial width $\Gamma(\Pi \to gg)$ of the $n = 1$, $l = 0$ ground state is given by

$$\Gamma(\Pi \to gg) = \frac{4}{3} M_\Pi N \alpha_S^2 \frac{|R(0)|^2}{M_\Pi^3}, \tag{10}$$

where the radial wavefunction at the origin for the ground state is

$$\frac{|R(0)|^2}{M_\Pi^3} = \frac{1}{16}\left[\frac{4}{3}\bar{\alpha}_s + C_N\bar{\alpha}_N + Q^2\bar{\alpha}\right]^3. \tag{11}$$

Considering again the example of $N = 2$ and $\bar{\alpha}_N = \bar{\alpha}_S = 0.15$, $\bar{\alpha} = 1/137$ we find

$$\sigma(pp \to \Pi)_{\rm DR} \approx 0.40 \text{ pb} \quad \text{at 13 TeV}. \tag{12}$$

Evidently, the direct resonance production cross section is indeed expected to be subdominant, around 8% that of the indirect production cross section [Eq. (5)].[4]

We now comment on the regime where Λ_N is smaller than $\Lambda_{\rm QCD}$. In fact, if the SU(N) confining scale is only a little smaller than $\Lambda_{\rm QCD}$ then a light quark pair can form out of the vacuum, leading to a bound state of two QCD color singlet states: $\chi\bar{q}$ and $\chi^\dagger q$. These color singlet states would themselves be bound together by SU(N) gauge interactions to form the SU(N) singlet bound state. Since only SU(N) interactions bind the two composite states ($\chi\bar{q}$ and $\chi^\dagger q$), it follows that $\frac{4}{3}\bar{\alpha}_s + C_N\bar{\alpha}_N + Q^2\bar{\alpha} \to C_N\bar{\alpha}_N + (Q - Q_q)^2\bar{\alpha}$ in Eqs. (8) and (11). Therefore if the confinement scale of SU(N) is smaller than that of QCD then the direct production rate becomes completely negligible relative to the indirect production mechanism. The rate of Π production is the same as that found earlier in Eq. (5), but the branching ratio to two photons is modified:

$$\text{Br}(\Pi \to \gamma\gamma) \simeq \frac{3NQ^4\alpha^2}{\frac{7}{3}N\alpha_S^2 + \frac{3}{2}C_N\alpha_N^2 + 3NQ^4\alpha^2}, \tag{13}$$

where, as before, we have neglected the small contribution of $\Pi \to Z\gamma/ZZ$ to the total width, and also the contribution from $\Pi \to hh$. In this regime somewhat larger values of Q can be accommodated, such as $Q = 5/6$ for $N = 2$.[5]

Notice that in the $\Lambda_N < \Lambda_{\rm QCD}$ regime the size of the mass splittings between the excited states becomes small as $\frac{4}{3}\bar{\alpha}_s + C_N\bar{\alpha}_N + Q^2\bar{\alpha} \to C_N\bar{\alpha}_N + (Q - Q_q)^2\bar{\alpha}$ in Eq. (8). We therefore expect no effective width enhancement due to the excited state decays at the LHC in the small Λ_N regime. Of course a larger effective width is still possible if there are several nearly degenerate scalar quirk states, which, as briefly mentioned earlier, can arise if χ transforms nontrivially under SU$(2)_L$.

IV. OTHER SIGNATURES

While the two photon decay channel of the bound state should be the most important signature, the dominant decay is expected to be via $\Pi \to gg$ and $\Pi \to \mathcal{H}\mathcal{H}$. The former process is expected to lead to dijet production while the latter will be an invisible decay. The dijet cross section is easily estimated:

$$\sigma(pp \to jj) \approx \begin{cases} 2.6N \times \text{Br}(\Pi \to gg) \text{ pb} & \text{at 13 TeV} \\ 0.5N \times \text{Br}(\Pi \to gg) \text{ pb} & \text{at 8 TeV} \end{cases}. \tag{14}$$

The limit from 8 TeV data is $\sigma(pp \to jj) \lesssim 2.5$ pb [29,30]. If gluons dominate the Π decays [i.e. $\text{Br}(\Pi \to gg) \approx 1$] then this experimental limit is satisfied for $N \leq 5$. For sufficiently large α_N the invisible decay can be enhanced, thereby reducing $\text{Br}(\Pi \to gg)$. In this circumstance the bound on N from dijet searches would weaken.

[3]Additional possibilities arise if χ transforms nontrivially under SU$(2)_L$, i.e. forming a representation $\mathbf{N}_L$. The mass degeneracy of the multiplet will be broken at tree-level by Higgs potential terms along with electroweak radiative corrections. The net effect is that the predicted width of the $pp \to \gamma\gamma$ bump can be effectively larger as there are N_L distinct bound states, Π^i, (of differing masses) which can each contribute to the decay width. Although each state is expected to have a narrow width, when smeared by the detector resolution the effect can potentially be a broad feature.

[4]If $\bar{\alpha}_N$ is sufficiently large, one can potentially have direct resonance production comparable or even dominating indirect production (such a scenario has been contemplated recently in [8,9]). Naturally at such large $\bar{\alpha}_N$ the perturbative calculations become unreliable, and one would have to resort to nonperturbative techniques such as lattice computations.

[5]Although it is perhaps too early to speculate on the possible role of χ in a more elaborate framework, we nevertheless remark here that particles fitting its description is required for spontaneous symmetry breaking of extended Pati-Salam type unified theories [28].

The invisible decays $\Pi \to \mathcal{H}\mathcal{H}$ are not expected to lead to an observable signal at leading order for much of the parameter space of interest.[6] However, the bremsstrahlung of a hard gluon from the initial state $pp \to \Pi g \to \mathcal{H}\mathcal{H}g$ can lead to a jet plus missing transverse energy signature. Current data are not expected to give stringent limits from such decay channels; however, this signature could become important when a larger data sample is collected. Note though that the rate will become negligible in the limit that α_N becomes small. Also, in the small Λ_N regime, where the bound state is formed from $\chi\bar{q}$ and $\chi^\dagger q$, the two-body decay $\Pi \to g\gamma$ (jet + photon) will also arise as in this case the scalar quirk pair is not necessarily in the color singlet configuration. The decay rate at leading order is substantial:

$$\frac{\Gamma(\Pi \to j\gamma)}{\Gamma(\Pi \to \gamma\gamma)} = \frac{8\alpha_s}{3\alpha Q^2}. \tag{15}$$

Nevertheless, we estimate that this is still consistent with current data [33], but would be expected to become important when a larger data sample is collected.

Another important signature of the model will be the $pp \to \Pi \to Z\gamma$ and $pp \to \Pi \to ZZ$ processes. The rates of these decays, relative to $\Pi \to \gamma\gamma$, are estimated to be

$$\frac{\Gamma(\Pi \to Z\gamma)}{\Gamma(\Pi \to \gamma\gamma)} = 2\tan^2\theta_W,$$
$$\frac{\Gamma(\Pi \to ZZ)}{\Gamma(\Pi \to \gamma\gamma)} = \tan^4\theta_W. \tag{16}$$

If χ transforms nontrivially under $SU(2)_L$ then deviations from these predicted rates arise along with the tree-level decay $\Pi \to W^+W^-$.

V. CONCLUSIONS

We have considered a charged scalar particle (χ) of mass around 375 GeV charged under both $SU(3)_C$ and a new confining gauge interaction [assigned to be SU(N) for definiteness]. These interactions confine $\chi^\dagger\chi$ into nonrelativistic bound states whose decays into photons can explain the 750 GeV diphoton excess observed at the LHC. Taking the new confining group to be SU(2), we found that the diphoton excess required χ to have electric charge approximately $Q \sim [\frac{1}{2}, 1]$. An important feature of our model is that the exotic particle χ has a mass much greater than the SU(N)-confinement scale Λ_N. In the absence of light SU(N)-charged matter fields this makes the dynamics of this new interaction qualitatively different to that of QCD: pair production of the scalars and the subsequent formation of the bound state dominates over direct bound state resonance production (at least in the perturbative regime where $\Lambda_N \lesssim \Lambda_{QCD}$). Since χ is a Lorentz scalar, decays of $\chi^\dagger\chi$ bound states to lepton pairs are naturally suppressed, and thus constraints from dilepton searches at the LHC can be ameliorated. This explanation is quite weakly constrained by current searches and data from the forthcoming run at the LHC will be able to probe our scenario more fully. In particular, dijet, monojet, di-Higgs, and jet + photon searches may be the most promising discovery channels.

ACKNOWLEDGMENTS

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[6] Scalar quirk loops can mediate hueball decays into gluons and other SM bosons [11,31,32]. The decay rate is uncertain, depending on the nonperturbative hueball dynamics. However, if the hueballs are able to decay within the detector then they can lead to observable signatures including displaced vertices. This represents another possible collider signature of the model.

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