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

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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 \cite{1,2}. The cross section $\sigma(pp \rightarrow \gamma\gamma)$ is estimated to be

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

and there is no evidence of any accompanying excess in the dilepton channel \cite{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 e.g. \cite{4} (for a recent discussion see also \cite{5}). The possibility that the new scalar particle is a bound state of exotic charged fermions has also been considered, e.g. \cite{6,11}. Here we consider the case that the 750 GeV state is a non-relativistic bound state constituted by an exotic scalar particle $\chi$ and its antiparticle, charged under SU(3)$_c$ as well as a new unbroken non-abelian gauge interaction. Having $\chi$ be a scalar rather than a fermion is not merely a matter of taste: In such a framework a fermionic $\chi$ 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 $\Pi$, can be produced through gluon–gluon fusion directly (i.e. at threshold $\sqrt{s} \sim M_\Pi$) or indirectly via $gg \rightarrow \chi\bar{\chi} \rightarrow \Pi + \text{soft quanta}$ (i.e. above $\Pi$ 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.

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, $\Lambda_\chi$. In fact we here consider only one such matter particle, $\chi$, so that $M_\chi \gg \Lambda_\chi$ is assumed. In this circumstance a $\chi\bar{\chi}$ 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 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 $\chi\bar{\chi}$ pair can be viewed as a highly excited bound state, which de-excites by SU($N$)-ball and soft glueball/pion emission \cite{11}.

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

$$\text{SU(3)$_c$} \otimes \text{SU(2)$_l$} \otimes U(1)$_Y$ \otimes \text{SU($N$)}.$$  \hfill (2)

This kind of theory can arise naturally in models which feature large colour groups \cite{12,14} and in models with leptonic colour \cite{15,17} but was also considered earlier by Okun \cite{18}. The notation \textit{quirks} for heavy particles charged under an unbroken gauge symmetry (where $M_\chi \gg \Lambda_\chi$) was introduced in \cite{11} where the relevant phenomenology was examined in some detail in a particular model \cite{1}. For convenience we borrow their nomenclature.

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\textsuperscript{†} Some other aspects of such models have been discussed over the
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The scalar χ that we introduce transforms under the extended gauge group (eq. 2) as

\[ \chi \sim (3, 1, Y; N), \]

where we use the normalization \( Q = Y/2 \). The possibility that \( \chi \) also transforms non-trivially under \( SU(2)_L \) is interesting, however for the purposes of this letter we focus on the \( SU(2)_L \) singlet case for definiteness. Since two-photon decays of non-relativistic quirkonium will be observed at the LHC, the mass of \( \chi \) will need to be around 375 GeV.

We have assumed that \( \chi \) 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 case where the \( SU(3)_C \) interaction strength and we have neglected the small contribution to the total width. Eq. 6 also neglects the decay to Higgs particles: \( \Pi \rightarrow hh \). 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 [22]. For a \( \chi \) mass of 375 GeV

\[ \sigma(pp \rightarrow \Pi \rightarrow \gamma\gamma) \sim 2.6 N \text{ pb at 13 TeV} \]

The branching fraction is to leading order:

\[ \text{Br}(\Pi \rightarrow \gamma\gamma) \approx \frac{3 N Q^4 \alpha^2}{3 N \alpha_s^2 + \frac{3}{2} C_N \alpha_s^2 + 3 N Q^4 \alpha_s^2}. \]

where \( C_N \equiv (N^2 - 1)/(2 N) \), and \( \alpha_N \) is the new \( SU(N) \) interaction strength and we have neglected the small contribution to the total width. Eq. 6 also neglects the decay to Higgs particles: \( \Pi \rightarrow hh \), which arises from the Higgs potential portal term \( \lambda_N \chi^\dagger \chi \phi \). Theoretically this rate is unconstrained given the dependence on the unknown parameter \( \lambda_N \), but could potentially be important. However, limits from resonant Higgs boson pair production derived from 13 TeV data: \( \sigma(pp \rightarrow X \rightarrow hh \rightarrow bbbb) \lesssim 50 \text{ fb at } M_X \approx 750 \text{ GeV} \) [23, 24] imply that the Higgs decay channel must indeed be subdominant (cf. \( \Pi \rightarrow gg, H \Pi \)).

The renormalized gauge coupling constants in eq. 6 are evaluated at the renormalization scale \( \mu \sim M_H/2 \). Taking for instance the specific case of \( N = 2 \), \( \alpha_N \approx \alpha_s \approx 0.10 \) (at \( \mu \sim M_H/2 \)) gives

\[ \sigma(pp \rightarrow \gamma\gamma) \approx 5 \left( \frac{Q}{1/2} \right)^4 \text{ fb at 13 TeV}. \]

At \( \sqrt{s} = 8 \text{ TeV} \) the cross section is around five times smaller. We present the cross section \( \sigma(pp \rightarrow \Pi \rightarrow \gamma\gamma) \) for a range of masses \( M_H \) and different combinations of \( Q \) and \( N \) in fig. 2. The parameter choice \( \alpha_N = \alpha_s \) and \( \Lambda_N = \Lambda_{\text{QCD}} \) has been assumed. (The cross section is not highly sensitive to \( \Lambda_N \), \( \alpha_N \) so long as we are in the perturbative

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and call the new quantum number \( hue \) and the massless gauge bosons \( huons (H) \).

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 \( \chi \) is a Lorentz scalar in light of previous work which indicated that bound states formed from a fermionic \( \chi \) 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][8]). The detailed consideration of a scalar \( \chi \) appears to have been largely overlooked, perhaps due to the paucity of known elementary scalar particles. With the recent discovery of a Higgs-like scalar at 125 GeV [20, 21] 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.
regime: \( \Lambda_N \lesssim \Lambda_{\text{QCD}} \). Evidently, for \( N = 2 \), a \( \chi \) with electric charge \( Q \approx 1/2 \) is produced at approximately the right rate to explain the diphoton excess.

In practice de-excitation 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 non-zero 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 de-excite predominately to \( l = 0 \) states rather than annihilate. Nevertheless, for sufficiently large \( \alpha_N \) 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:

\[
E_n = -\frac{1}{8n^2}\left[\frac{4}{3}\bar{\alpha}_s + C_N\bar{\alpha}_N + Q^2\bar{\alpha}\right]^2.
\]

The above formula was adapted from known results with quarkonium, e.g. [25] (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: \( a_0 = 4/[4\bar{\alpha}_s/3 + C_N\bar{\alpha}_N + Q^2\bar{\alpha}]M_H \). The bound state, described by the radial quantum number \( n \) has mass given by \( M_H(n) = 2M_X + E_n \). Considering 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 \( \Delta M = (E_1 - E_2) \approx 0.01M_H \). Larger mass splittings will be possible if \( \bar{\alpha}_s > \bar{\alpha}_N \), although it has been shown in the context of fermionic quirk models that the phenomenology is substantially altered in this regime [27]. In particular, the huebells can become so heavy that the decays of the bound state into huebells is kinematically forbidden.

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

\[
\sigma(pp \to \Pi)_{\text{dir}} \approx \frac{C_{gg}K_{gg}\Gamma(\Pi \to gg)}{sM_H},
\]

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 \( C_{gg} \approx 2137 \) and \( K_{gg} = 1.6 \) [28]. The partial width \( \Gamma(\Pi \to gg) \) of the \( n = 1, l = 0 \) ground state is given by

\[
\Gamma(\Pi \to gg) = \frac{4}{3}M_HN\alpha_s^2|R(0)|^2M_H^3,
\]

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

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

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

\[
\sigma(pp \to \Pi)_{\text{dir}} \approx 0.40 \text{ pb} \quad \text{at 13 TeV}.
\]

Evidently, the direct resonance production cross section is indeed expected to be subdominant, around 8% that 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 \( N \) distinct bound states, \( \Pi^0 \), (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.

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\[3\] Additional possibilities arise if \( \chi \) transforms nontrivially under SU(2), i.e. forming a representation \( \mathbf{N}_L \). The mass degeneracy of the multiplet will be broken at tree-level by Higgs potential.
FIG. 2. The cross section $\sigma(pp \rightarrow \Pi \rightarrow \gamma\gamma)$ at 13 TeV for a range of quirkonium masses $M_\Pi$ 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 errors associated with the parton distribution functions.

We now comment on the regime where $\Lambda_n$ is smaller than $\Lambda_{QCD}$. In fact, if the SU($N$) confining scale is only a little smaller than $\Lambda_{QCD}$ 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 $\frac{4}{3} \tilde{\alpha}_s + C_n \tilde{\alpha}_n + Q^2 \tilde{\alpha} \rightarrow C_n \tilde{\alpha}_n + (Q - Q_q)^2 \tilde{\alpha}$ 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 $\Pi$ production is the same as that found earlier in eq. 5 but the branching ratio to two photons is modified:

$$\text{Br}(\Pi \rightarrow \gamma\gamma) \simeq \frac{3NQ^4\alpha^2}{\frac{2}{3}N\alpha_s^2 + \frac{3}{2}C_n\alpha_n^2 + 3NQ^4\alpha^2}$$

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

Notice that in the $\Lambda_n < \Lambda_{QCD}$ regime the size of the mass splittings between the excited states becomes small as $\frac{4}{3} \tilde{\alpha}_s + C_n \tilde{\alpha}_n + Q^2 \tilde{\alpha} \rightarrow C_n \tilde{\alpha}_n + (Q - Q_q)^2 \tilde{\alpha}$ in eq. 8.

We therefore expect no effective width enhancement due to the excited state decays at the LHC in the small $\Lambda_n$ 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 $\chi$ transforms nontrivially under SU(2)$_L$.

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 $\Pi \rightarrow gg$ and $\Pi \rightarrow HH$. The former process is expected to lead to dijet production while the latter will be an invisible decay. The dijet cross

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4 If $\tilde{\alpha}_n$ 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 $\tilde{\alpha}_n$ the perturbative calculations become unreliable, and one would have to resort to non-perturbative techniques such as lattice computations.

5 Although it is perhaps too early to speculate on the possible role of $\chi$ in a more elaborate framework, we nevertheless remark here that particles fitting its description are required for spontaneous symmetry breaking of extended Pati-Salam type unified theories [27].
section is easily estimated:

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

(14)

The limit from 8 TeV data is $\sigma(pp \rightarrow jj) \lesssim 2.5$ pb [28, 29]. If gluons dominate the $\Pi$ decays (i.e. $\text{Br}(\Pi \rightarrow gg) \approx 1$) then this experimental limit is satisfied for $N \leq 5$. For sufficiently large $\alpha_s$ the invisible decay can be enhanced, thereby reducing $\text{Br}(\Pi \rightarrow gg)$. In this circumstance the bound on $N$ from dijet searches would weaken.

The invisible decays $\Pi \rightarrow \mathcal{HH}$ 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: $pp \rightarrow \Pi g \rightarrow \mathcal{HH}g$ 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 $\alpha_s$ becomes small. Also, in the small $\Lambda_R$ regime, where the bound state is formed from $\chi \bar{q}$ and $\chi' \bar{q}$, the two-body decay $\Pi \rightarrow g\gamma$ (jet + photon) will also arise as in this case the scalar quark pair is not necessarily in the color singlet configuration. The decay rate at leading order is substantial:

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

Nevertheless, we estimate that this is still consistent with current data [30], but would be expected to become important when a larger data sample is collected.

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

$$\frac{\Gamma(\Pi \rightarrow Z\gamma)}{\Gamma(\Pi \rightarrow \gamma\gamma)} = 2\tan^2\theta_W, \quad (16)$$

$$\frac{\Gamma(\Pi \rightarrow ZZ)}{\Gamma(\Pi \rightarrow \gamma\gamma)} = \tan^4\theta_W.$$  

CONCLUSIONS

We have considered a charged scalar particle $\chi$ 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 non-relativistic 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 $\chi$ to have electric charge approximately $Q \sim \{\frac{1}{2}, 1\}$. An important feature of our model is that the exotic particle $\chi$ has a mass much greater than the SU($N$)-confinement scale $\Lambda_N$. 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 $\Lambda_N \lesssim \Lambda_{\text{QCD}}$). Since $\chi$ 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, mono-jet, di-Higgs and jet + photon searches may be the most promising discovery channels.

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Scalar quark loops can mediate hueball decays into gluons and other SM bosons [11, 33, 34]. The decay rate is uncertain, depending on the non-perturbative 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.

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