

A composite model for the 750 GeV diphoton excess

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ABSTRACT: We study a simple model in which the recently reported 750 GeV diphoton excess arises from a composite pseudo Nambu-Goldstone boson — hidden pion — produced by gluon fusion and decaying into two photons. The model only introduces an extra hidden gauge group at the TeV scale with a vectorlike quark in the bifundamental representation of the hidden and standard model gauge groups. We calculate the masses of all the hidden pions and analyze their experimental signatures and constraints. We find that two colored hidden pions must be near the current experimental limits, and hence are probed in the near future. We study physics of would-be stable particles — the composite states that do not decay purely by the hidden and standard model gauge dynamics — in detail, including constraints from cosmology. We discuss possible theoretical structures above the TeV scale, e.g. conformal dynamics and supersymmetry, and their phenomenological implications. We also discuss an extension of the minimal model in which there is an extra hidden quark that is singlet under the standard model and has a mass smaller than the hidden dynamical scale. This provides two standard model singlet hidden pions that can both be viewed as diphoton/diboson resonances produced by gluon fusion. We discuss several scenarios in which these (and other) resonances can be used to explain various excesses seen in the LHC data.

KEYWORDS: Beyond Standard Model, Technicolor and Composite Models

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Contents

1	Introduction	1
2	Model at the TeV scale	3
2.1	Hidden pion for the 750 GeV diphoton excess	3
2.2	Other hidden pions	5
2.3	Hidden eta prime	8
2.4	Heavy spin-1 resonances	9
3	Physics at higher energies	10
3.1	Physics of hidden pion χ	10
3.2	Cosmology	13
3.3	Possible ultraviolet theories	17
4	An extra hidden light quark	19
4.1	Diphoton (diboson) signals and other phenomenology	20
4.2	Hidden quasi-conserved quantities and cosmology	24
5	Discussion	25
A	Hidden sector CP	26

1 Introduction

If the recently announced diphoton excess at $\simeq 750$ GeV [1–4] remains as a true signal, it indicates a long-awaited discovery of new physics beyond the standard model. In a recent paper [5], we have proposed that this excess may result from a composite spin-0 particle decaying into a two-photon final state. Among the possibilities discussed, here we concentrate on the case in which the particle is a composite pseudo Nambu-Goldstone boson associated with new strong dynamics at the TeV scale, which is singly produced by gluon fusion and decays into two photons. For works that have appeared around the same time and discussed similar models to those in [5], see [6–9]; other related works include [10–15]. A class of theories involving similar dynamics with vectorlike matter charged under both hidden and standard model gauge groups was studied in [16, 17]. The possibility of obtaining standard model dibosons from a composite scalar particle was utilized in a different context in [18–20].

In this paper, we study a simple model presented in ref. [5] which introduces an extra gauge group $G_H = \text{SU}(N)$ at the TeV scale, in addition to the standard model gauge group $G_{\text{SM}} = \text{SU}(3)_C \times \text{SU}(2)_L \times \text{U}(1)_Y$, with extra matter — hidden quarks — in the vectorlike bifundamental representation of G_H and $\text{SU}(5) \supset G_{\text{SM}}$. Here, $\text{SU}(5)$ is used only as a

simple mnemonic; it does not mean that the three factors of G_{SM} are unified at the TeV scale. In this model, the diphoton resonance is one of the pseudo Nambu-Goldstone bosons — hidden pions — which is neutral under G_{SM} and has the mass of $\simeq 750$ GeV. We consider this particular model here because it is theoretically simple and leads to predictions that can be tested in the near future. The only free parameters of the model, beyond the size N of G_H , are the dynamical scale of G_H and two (in general complex) masses for the hidden quarks, out of which two numbers are fixed by the mass and diphoton rate of the resonance. We calculate all the masses of the hidden pions, which for a given N depend only on a single free parameter: the ratio of the absolute values of the two hidden quark masses. We find that two colored hidden pions must be near the current experimental limits (unless N is unreasonably large); they respectively lead to narrow dijet, Z -jet, and γ -jet resonances below ~ 1.6 TeV and to heavy stable charged and neutral hadrons (or leptoquark-type resonances) below ~ 1.2 TeV. We also discuss other, higher resonances in the model, which are expected to be in the multi-TeV region, as well as the effect of possible CP violation in the G_H sector on the standard model physics. Furthermore, we investigate what the structure of the model above the TeV region can be. This includes the possibility of (a part of) the theory being conformal and/or supersymmetric. This affects collider signatures and cosmological implications of one of the colored hidden pions: the one that does not decay through G_H or standard model gauge dynamics. We also perform detailed analyses of cosmology of hidden baryons.

We then discuss an extension of the minimal model which has an extra hidden quark that is singlet under the standard model gauge group. One of the salient features of this model is that there are two diphoton (diboson) resonances in the hidden pion sector, which can both be produced by gluon fusion and decay into two electroweak gauge bosons. We consider several scenarios associated with these and other related resonances. Representative ones are: (i) the two resonances correspond to two diphoton “excesses” seen in the ATLAS data at $\simeq 750$ GeV and $\simeq 1.6$ TeV [1, 2] (although the latter is much weaker than the former); (ii) the two resonances are both near $\simeq 750$ GeV with a mass difference of 10s of GeV, explaining a slight preference to a wide width in the ATLAS data; (iii) the two resonances are at $\simeq 750$ GeV and $\simeq 2$ TeV, responsible for the 750 GeV diphoton excess [1–4] and the 2 TeV diboson excess [21], respectively. We calculate the masses of the hidden pions in the model and discuss their phenomenology. We find that the masses of the hidden pions can be larger than the case without the singlet hidden quark; in particular, the leptoquark type hidden pion can be as heavy as ~ 1.5 TeV, depending on scenarios. We discuss physics of hidden pions and hidden baryons that decay only through interactions beyond the G_H and standard model gauge dynamics. We find that cosmological constraints on this model are weaker than those in the model without the extra hidden quark.

The organization of this paper is as follows. In section 2, we consider the minimal model and its phenomenology at the TeV scale. We calculate all the hidden pion masses and discuss their signatures and current constraints. We also discuss particles with higher masses, in particular the hidden η' meson and spin-1 resonances. In section 3, we study physics above the TeV scale, especially its implications for collider physics and cosmology. The hidden pion that is stable under the G_H and G_{SM} gauge dynamics as well as low-lying

	$G_H = \text{SU}(N)$	$\text{SU}(3)_C$	$\text{SU}(2)_L$	$\text{U}(1)_Y$
Ψ_D	\square	$\bar{\square}$	$\mathbf{1}$	$1/3$
Ψ_L	\square	$\mathbf{1}$	\square	$-1/2$
$\bar{\Psi}_D$	$\bar{\square}$	\square	$\mathbf{1}$	$-1/3$
$\bar{\Psi}_L$	$\bar{\square}$	$\mathbf{1}$	\square	$1/2$

Table 1. Charge assignment of hidden quarks under the hidden and standard model gauge groups. Here, $\Psi_{D,L}$ and $\bar{\Psi}_{D,L}$ are left-handed Weyl spinors.

hidden baryons are studied in detail. We discuss the possibility that the G_H sector is conformal and/or that the theory is supersymmetric above the TeV scale. In section 4, we study an extension of the model in which there is an extra hidden quark that is singlet under G_{SM} and has a mass smaller than Λ . We discuss possible signals of two G_{SM} -singlet hidden pions which can be viewed as diboson resonances produced by gluon fusion. Section 5 is devoted to final discussion. In the appendix, we analyze the effect of possible CP violation in the G_H sector on the standard model physics.

2 Model at the TeV scale

The model at the TeV scale is given by a hidden gauge group $G_H = \text{SU}(N)$, with the dynamical scale (the mass scale of generic low-lying resonances) Λ , and hidden quarks charged under both G_H and the standard model gauge groups as shown in table 1.¹ The hidden quarks have mass terms

$$\mathcal{L} = -m_D \Psi_D \bar{\Psi}_D - m_L \Psi_L \bar{\Psi}_L + \text{h.c.}, \tag{2.1}$$

where we take $m_{D,L} > 0$, which does not lead to a loss of generality if we keep all the phases in the other part of the theory. These masses are assumed to be sufficiently smaller than the dynamical scale, $m_{D,L} \ll \Lambda$, so that $\Psi_{D,L}$ and $\bar{\Psi}_{D,L}$ can be regarded as light quarks from the point of view of the G_H dynamics. Note that the charge assignment of the hidden quarks is such that they are a vectorlike fermion in the bifundamental representation of G_H and $\text{SU}(5) \supset G_{\text{SM}}$. The model therefore preserves gauge coupling unification at the level of the standard model; this is significant especially given the possible threshold corrections around the TeV and unification scales (see, e.g., [22]). The unification of the couplings becomes even better if we introduce supersymmetry near the TeV scale (see section 3).

2.1 Hidden pion for the 750 GeV diphoton excess

The strong G_H dynamics makes the hidden quarks condensate

$$\langle \Psi_D \bar{\Psi}_D + \Psi_D^\dagger \bar{\Psi}_D^\dagger \rangle \approx \langle \Psi_L \bar{\Psi}_L + \Psi_L^\dagger \bar{\Psi}_L^\dagger \rangle \equiv -c. \tag{2.2}$$

¹Throughout the paper, we adopt the hypercharge normalization such that the standard model left-handed Weyl fermions have $(q, u, d, l, e) = (1/6, -2/3, 1/3, -1/2, 1)$.

These condensations do not break the standard model gauge groups, since the hidden quark quantum numbers under these gauge groups are vectorlike with respect to G_H [23]. The spectrum below Λ then consists of hidden pions, arising from spontaneous breaking of the approximate $SU(5)_A$ axial flavor symmetry:

$$\psi(\mathbf{Adj}, \mathbf{1}, 0), \quad \chi\left(\square, \square, -\frac{5}{6}\right), \quad \varphi(\mathbf{1}, \mathbf{Adj}, 0), \quad \phi(\mathbf{1}, \mathbf{1}, 0), \quad (2.3)$$

where ψ , φ , and ϕ are real scalars while χ is a complex scalar, and the quantum numbers represent those under $SU(3)_C \times SU(2)_L \times U(1)_Y$. The masses of these particles are given by [24]

$$m_\psi^2 = 2m_D \frac{c}{f^2} + 3\Delta_C, \quad (2.4)$$

$$m_\chi^2 = (m_D + m_L) \frac{c}{f^2} + \frac{4}{3}\Delta_C + \frac{3}{4}\Delta_L + \frac{5}{12}\Delta_Y, \quad (2.5)$$

$$m_\varphi^2 = 2m_L \frac{c}{f^2} + 2\Delta_L, \quad (2.6)$$

$$m_\phi^2 = \frac{4m_D + 6m_L}{5} \frac{c}{f^2}. \quad (2.7)$$

Here, f is the decay constant,² and $\Delta_{C,L,Y}$ are contributions from standard model gauge loops, of order

$$\Delta_C \simeq \frac{3g_3^2}{16\pi^2}\Lambda^2, \quad \Delta_L \simeq \frac{3g_2^2}{16\pi^2}\Lambda^2, \quad \Delta_Y \simeq \frac{3g_1^2}{16\pi^2}\Lambda^2, \quad (2.8)$$

where g_3 , g_2 , and g_1 are the gauge couplings of $SU(3)_C$, $SU(2)_L$, and $U(1)_Y$, respectively, with g_1 in the $SU(5)$ normalization. Using naive dimensional analysis [25, 26], we can estimate the quark bilinear condensate and the decay constant as

$$c \simeq \frac{N}{16\pi^2}\Lambda^3, \quad f \simeq \frac{\sqrt{N}}{4\pi}\Lambda, \quad (2.9)$$

where we have assumed $N \gtrsim 5$, i.e. the number of color is not much smaller than that of flavor in the G_H gauge theory. For $N \ll 5$, we might instead have $c \approx (5/16\pi^2)\Lambda^3$ and $f \approx (\sqrt{5}/4\pi)\Lambda$, but below we use eq. (2.9) even in this case because the resulting differences are insignificant for our results.

The couplings of the hidden pions with the standard model gauge fields are determined by chiral anomalies and given by

$$\begin{aligned} \mathcal{L} = & \frac{Ng_3^2}{64\pi^2 f} d^{abc} \psi^a \epsilon^{\mu\nu\rho\sigma} G_{\mu\nu}^b G_{\rho\sigma}^c + \frac{Ng_3 g_1}{32\sqrt{15}\pi^2 f} \psi^a \epsilon^{\mu\nu\rho\sigma} G_{\mu\nu}^a B_{\rho\sigma} \\ & - \frac{3Ng_2 g_1}{64\sqrt{15}\pi^2 f} \varphi^\alpha \epsilon^{\mu\nu\rho\sigma} W_{\mu\nu}^\alpha B_{\rho\sigma} + \frac{Ng_3^2}{32\sqrt{15}\pi^2 f} \phi \epsilon^{\mu\nu\rho\sigma} G_{\mu\nu}^a G_{\rho\sigma}^a \\ & - \frac{3Ng_2^2}{64\sqrt{15}\pi^2 f} \phi \epsilon^{\mu\nu\rho\sigma} W_{\mu\nu}^\alpha W_{\rho\sigma}^\alpha - \frac{Ng_1^2}{64\sqrt{15}\pi^2 f} \phi \epsilon^{\mu\nu\rho\sigma} B_{\mu\nu} B_{\rho\sigma}, \end{aligned} \quad (2.10)$$

²Our definition of the decay constant, f , is a factor of 2 different from that in ref. [24]: $f = F/2$.

where $a, b, c = 1, \dots, 8$ and $\alpha = 1, 2, 3$ are $SU(3)_C$ and $SU(2)_L$ adjoint indices, respectively, and $d^{abc} \equiv 2\text{tr}[t^a\{t^b, t^c\}]$ with t^a being half of the Gell-Mann matrices. We assume that the ϕ particle produced by gluon fusion and decaying to a diphoton is responsible for the observed excess [5], so we take

$$m_\phi \simeq 750 \text{ GeV}. \quad (2.11)$$

The decay of ϕ occurs through interactions in eq. (2.10) and leads to standard model gauge bosons. The diphoton rate at $\sqrt{s} = 13 \text{ TeV}$ is given by

$$\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma) \simeq 7.8 \text{ fb} \left(\frac{N 500 \text{ GeV}}{5 f} \right)^2. \quad (2.12)$$

Here, we have used the NNPDF 3.0 parton distribution function [27] and determined the overall normalization (the QCD K factor) such that it reproduces the production cross section of a standard model-like Higgs boson of mass 750 GeV at $\sqrt{s} = 14 \text{ TeV}$ [28]. Since the observed excess corresponds to $\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma) \simeq (6 \pm 2) \text{ fb}$ [29, 30], this gives

$$f \simeq 570 \text{ GeV} \frac{N}{5} \sqrt{\frac{6 \text{ fb}}{\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma)}}. \quad (2.13)$$

With this value of f , the upper limits from searches in the 8 TeV data [31, 32] are evaded.

The ratios of branching fractions to various ϕ decay modes are given by

$$\begin{aligned} \frac{B_{\phi \rightarrow gg}}{B_{\phi \rightarrow \gamma\gamma}} &= 8 \left(\frac{6g_3^2}{14e^2} \right)^2 \simeq 200, & \frac{B_{\phi \rightarrow WW}}{B_{\phi \rightarrow \gamma\gamma}} &= 2 \left(\frac{9}{14 \sin^2 \theta_W} \right)^2 \simeq 15, \\ \frac{B_{\phi \rightarrow ZZ}}{B_{\phi \rightarrow \gamma\gamma}} &= \left(\frac{9 + 5 \tan^4 \theta_W}{14 \tan^2 \theta_W} \right)^2 \simeq 5, & \frac{B_{\phi \rightarrow Z\gamma}}{B_{\phi \rightarrow \gamma\gamma}} &= 2 \left(\frac{9 - 5 \tan^2 \theta_W}{14 \tan \theta_W} \right)^2 \simeq 2, \end{aligned} \quad (2.14)$$

where e and θ_W are the electromagnetic coupling and the Weinberg angle, respectively, and we have ignored the phase space factors. These values are consistent with the constraints from searches of high-mass diboson and dijet resonances in the 8 TeV data [5, 7, 33]. Observing these decay modes in the 13 TeV run with the predicted rates would provide an important test of the model.

2.2 Other hidden pions

The identification of ϕ as the 750 GeV diphoton resonance leads, through eq. (2.7), to

$$\frac{2m_D + 3m_L}{5} \sim 90 \text{ GeV} \sqrt{\frac{5}{N}} \sqrt{\frac{\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma)}{6 \text{ fb}}}, \quad (2.15)$$

where we have used eqs. (2.9), (2.13). This, however, leaves the ratio $r \equiv m_D/m_L$ undetermined. With $\Lambda = 3.2 \text{ TeV} \sqrt{N/5}$, which is motivated by eqs. (2.9), (2.13), the masses of the other hidden pions are determined by eqs. (2.4)–(2.6) in terms of r :

$$m_\psi^2 \simeq \frac{5r}{2r+3} (750 \text{ GeV})^2 + \frac{N}{5} (760 \text{ GeV})^2, \quad (2.16)$$

$$m_\chi^2 \simeq \frac{5r+5}{4r+6} (750 \text{ GeV})^2 + \frac{N}{5} (580 \text{ GeV})^2, \quad (2.17)$$

$$m_\varphi^2 \simeq \frac{5}{2r+3} (750 \text{ GeV})^2 + \frac{N}{5} (400 \text{ GeV})^2. \quad (2.18)$$

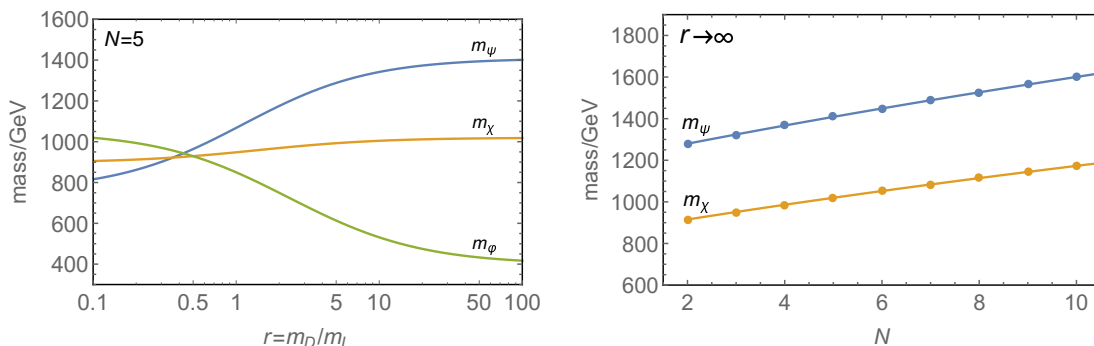


Figure 1. The masses of hidden pions ψ , χ , and φ for $m_\phi = 750$ GeV as functions of $r = m_D/m_L$ for $N = 5$ (left) and the maximal values of the ψ and χ masses, $m_\psi|_{r \rightarrow \infty}$ and $m_\chi|_{r \rightarrow \infty}$, as functions of N (right). Here, we have taken $\Lambda = 3.2 \text{ TeV} \sqrt{N/5}$, motivated by eqs. (2.9), (2.13), and used eq. (2.8) with unit coefficients.

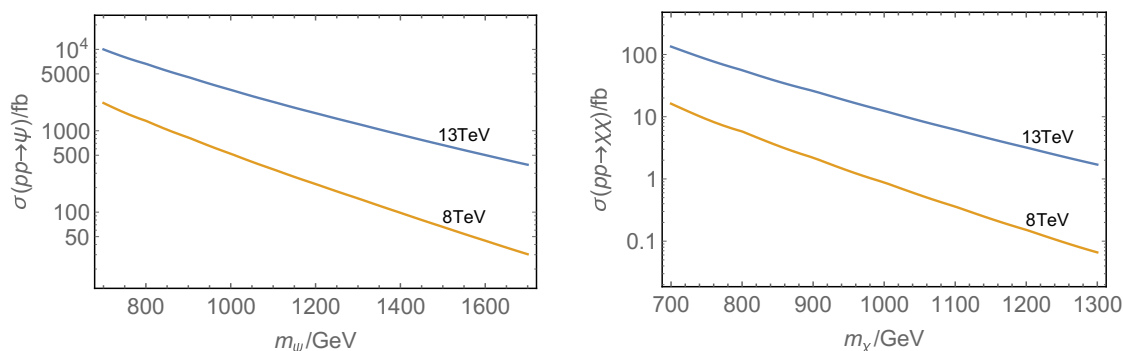


Figure 2. Cross section of single production of ψ through gluon fusion assuming $\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma) = 6 \text{ fb}$ (left) and that of pair production of χ (right). Both are plotted for $\sqrt{s} = 13 \text{ TeV}$ and 8 TeV .

Here, in the second terms we have used eq. (2.8) with unit coefficients, but we do not expect that using the true coefficients (which are not known in general) make a significant difference. In the left panel of figure 1, we plot these masses for $N = 5$ as functions of r . In the right panel, we plot the maximal values of the ψ and χ masses, $m_\psi|_{r \rightarrow \infty}$ and $m_\chi|_{r \rightarrow \infty}$, as functions of N . We find that these particles are at

$$m_\psi \lesssim 1.6 \text{ TeV}, \quad m_\chi \lesssim 1.2 \text{ TeV}, \quad (2.19)$$

unless N is very large, $N \geq 10$. We note that if m_D and m_L are unified at a conventional unification scale (around 10^{14} – 10^{17} GeV), then their ratio at the TeV scale is in the range $r \simeq 1.5 - 3$, with the precise value depending on the structure of the theory above the TeV scale.

The ψ particle is created dominantly through single production by gluon fusion, whose cross section is plotted in the left panel of figure 2 for $\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma) = 6 \text{ fb}$. (The dependencies on f and N cancel for a fixed value of $\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma)$.) Once produced, ψ

decays via interactions in eq. (2.10) to $g\gamma$, gZ , and gg with the branching fractions

$$B_{\psi \rightarrow g\gamma} = \frac{2g_1^2 \cos^2\theta_W}{25g_3^2 + 2g_1^2} \simeq 0.01, \quad B_{\psi \rightarrow gZ} = \frac{2g_1^2 \sin^2\theta_W}{25g_3^2 + 2g_1^2} \simeq 0.003, \quad (2.20)$$

and $B_{\psi \rightarrow gg} = 1 - B_{\psi \rightarrow gZ} - B_{\psi \rightarrow g\gamma}$. The lower bound on the mass of ψ from the LHC data so far [34–37] is about 800 GeV. For $N < 5$, this requires r to be not significantly smaller than 1.

The χ particle is produced only in pairs because of the conservation of “ D number” and “ L number” at the renormalizable level, under which Ψ_D and Ψ_L transform as $(1, 0)$ and $(0, 1)$, respectively. The production cross section is plotted in the right panel of figure 2, ignoring possible form factors which may become important when the hierarchy between m_χ and Λ is not significant due to small N , e.g. $N = 3$. The signal of χ depends on its lifetime, which is determined by the strength of nonrenormalizable interactions between the hidden quarks and the standard model particles violating D and L numbers. (This issue will be discussed in section 3.1.) Consider first the case in which χ is stable at collider timescales. In this case, a produced χ particle picks up a light quark of the standard model, becoming a heavy fermionic “hadron.” Since the charge $\pm 4/3$ component of χ is heavier than the charge $\pm 1/3$ component by about 700 MeV [38], the former is subject to weak decays into the latter with $c\tau \lesssim 1$ mm, so that the final heavy hadron has a charge 0 or ± 1 . The mass splitting between these neutral and charged hadrons is of order MeV, so that the weak decay between them is slow and they can both be regarded as stable particles for usual collider purposes. The lower bound on the χ mass can thus be read off from that of the stable bottom squark [39, 40] with the doubled production cross section as given in figure 2 (because of the twice larger number of degrees of freedom). The bound is about 900 GeV. On the other hand, if the nonrenormalizable interactions are strong, χ may decay promptly into a quark and a lepton. In this case, the lower bound on the χ mass is about 750 GeV–1 TeV, depending on the details of the decay modes [41, 42]. In any event, because of the theoretical expectations in eq. (2.19), searches of the ψ and χ particles provide important probes of the model.

Finally, the φ particle is standard model color singlet, so it can be produced only through electroweak processes or decays of heavier resonances. The decay of φ occurs through interactions of eq. (2.10). The φ^\pm decays into $W\gamma$ and WZ with the branching fractions of $\cos^2\theta_W \simeq 0.77$ and $\sin^2\theta_W \simeq 0.23$, respectively, while φ^0 decays into $\gamma\gamma$, γZ and ZZ with the branching fractions $\sin^2(2\theta_W)/2 \simeq 0.35$, $\cos^2(2\theta_W) \simeq 0.30$, and $\sin^2(2\theta_W)/2 \simeq 0.35$, respectively. The current bounds on $\varphi^{\pm,0}$ are weak and do not constrain the model further.

After electroweak symmetry breaking, ϕ and φ can mix via the operator $\mathcal{L} \sim (g_2^4/16\pi^2)(m_L/\Lambda)\phi\varphi^\alpha h^\dagger\sigma^\alpha h$, which is generated by a loop involving electroweak gauge bosons. Here, h is the standard model Higgs field, and the factor of m_L appears because the operator breaks the shift symmetry of ϕ . The resultant mixing angle, however, is small, $\theta_{\phi\varphi} \sim (g_2^4/16\pi^2)(m_L\langle h \rangle^2/m_\phi^2\Lambda) \sim 10^{-6}$, so that it does not affect the phenomenology of hidden pions or electroweak symmetry breaking.

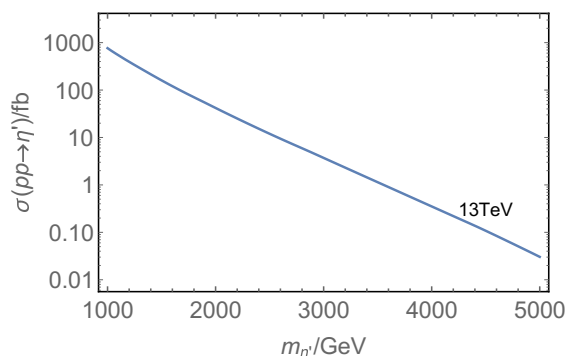


Figure 3. Single production cross section of the hidden η' through gluon fusion at $\sqrt{s} = 13$ TeV, calculated using eq. (2.22) with N/f determined by eq. (2.13) with $\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma) = 6$ fb.

2.3 Hidden eta prime

Another interesting particle arising from the G_H sector is the hidden η' state associated with the $U(1)_A$ axial symmetry, which we simply call η' below. The mass of this particle is expected to be at the dynamical scale [43, 44]

$$m_{\eta'} \approx \sqrt{\frac{5}{N}} \Lambda \approx 3.2 \text{ TeV} \sqrt{\frac{6 \text{ fb}}{\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma)}}, \quad (2.21)$$

where we have used eqs. (2.9), (2.13) in the last expression. The couplings to the standard model gauge bosons can be estimated by $U(1)_A$ anomalies with respect to the standard model gauge groups

$$\mathcal{L} \approx \frac{N g_3^2}{32\sqrt{10}\pi^2 f} \eta' \epsilon^{\mu\nu\rho\sigma} G_{\mu\nu}^a G_{\rho\sigma}^a + \frac{N g_2^2}{32\sqrt{10}\pi^2 f} \eta' \epsilon^{\mu\nu\rho\sigma} W_{\mu\nu}^\alpha W_{\rho\sigma}^\alpha + \frac{N g_1^2}{32\sqrt{10}\pi^2 f} \eta' \epsilon^{\mu\nu\rho\sigma} B_{\mu\nu} B_{\rho\sigma}. \quad (2.22)$$

This expression is valid in the large N limit, and we expect that it gives a good approximation even for moderately large N . The production cross section through gluon fusion calculated using this expression is depicted in figure 3. The energy dependence of the QCD K factor is at most of $O(10\%)$ and hence is neglected.

Dominant decay modes of η' depend strongly on whether the G_H sector respects CP (or parity) or not. In general, there is the possibility that the G_H sector has significant CP violation due to complex phases of the hidden quark masses or the θ parameter for the G_H gauge theory. This possibility is particularly natural if there is a QCD axion, eliminating the effect of CP violation on the QCD θ parameter (but no axion acting on the G_H sector); see the appendix. In this case, η' decays into two hidden pions through CP violating terms of the form

$$\mathcal{L} \sim \frac{mc}{f^3} \eta' \pi \pi, \quad (2.23)$$

where m and π collectively denote hidden quark masses and hidden pion fields, respectively. The final state then consists of either 2 quasi-stable particles, $\chi\chi$, or 4 standard model gauge bosons. As discussed in the appendix, CP violation of the G_H sector may also be observed in the neutron electric dipole moment.

If the G_H sector respects CP , e.g. as in the case in which the axion mechanism also operates in the G_H sector, then η' decays dominantly into 3π or 2 standard model gauge bosons. If the former is kinematically open, it dominates the decay; this will be the case if the η' mass is indeed given by eq. (2.21) with the coefficient close to unity. On the other hand, if the 3π mode is kinematically forbidden due to a (somewhat unexpected) suppression of the coefficient of eq. (2.21), then the decay is to 2 standard model gauge bosons. Since the mixing between η' and hidden pions is suppressed due to small hidden quark masses, the branching ratios of η' are determined purely by the $SU(5)$ flavor symmetry and given by

$$R_{gg} : R_{WW} : R_{ZZ} : R_{Z\gamma} : R_{\gamma\gamma} \approx \frac{8g_3^4}{g_2^4} : 2 : \left(c_W^2 + \frac{5s_W^4}{3c_W^2} \right)^2 : 2 \left(s_W c_W - \frac{5s_W^3}{3c_W} \right)^2 : \frac{64s_W^4}{9} \\ \simeq 0.94 : 0.038 : 0.015 : 0.0017 : 0.0072. \quad (2.24)$$

Here, we have used the abbreviations $R_{AB} = B_{\eta' \rightarrow AB}$, $s_W = \sin \theta_W$, and $c_W = \cos \theta_W$. If these modes dominate, the production and decay of η' also leads to a diphoton signal.³

2.4 Heavy spin-1 resonances

We finally discuss spin-1 resonances in the G_H sector having odd C and P , which we refer to as hidden ρ mesons. We denote the hidden ρ mesons that have the same flavor $SU(5)$ charges as ψ , χ , φ , ϕ , and η' by ρ_ψ , ρ_χ , ρ_φ , ρ_ϕ , and ω , respectively. These hidden ρ mesons are expected to be as heavy as Λ . They are produced at the LHC and yield interesting signals. For earlier studies of phenomenology of spin-1 resonant production of pairs of hidden particles, see [16, 17].

The ρ_ψ particle mixes with the standard model gluon and couples to standard model quarks with a coupling constant $\sim \sqrt{N}g_3^2/4\pi$. The single production cross section of ρ_ψ from quark initial states is of order 1000 – 1 fb for the ρ_ψ mass of 2 – 5 TeV at the 13 TeV LHC. It is also singly produced from a two gluon initial state via higher dimensional operators suppressed by Λ

$$\mathcal{L} \approx i \frac{\sqrt{N}g_3^2}{4\pi\Lambda^2} f^{abc} (D_\mu \rho_\nu^a - D_\nu \rho_\mu^a) G_\rho^{b\mu} G^{c\rho\nu}. \quad (2.25)$$

The production cross section from the two gluon initial state is roughly comparable to that from quark initial states. The produced ρ_ψ mainly decays into a pair of ψ with a large width, if it is kinematically allowed; the resulting ψ in turn decays into gg , gZ , or $g\gamma$, yielding a narrow dijet, Z -jet, or γ -jet resonance. The decay of ρ_ψ into $\psi\phi$ is forbidden by C -parity.

³It is amusing to identify this as the origin of the slight “excess” at $\simeq 1.6$ TeV in the ATLAS diphoton data [1, 2] (although this requires a deviation from the naive estimate of the η' mass, eq. (2.21), by a factor of 2). Indeed, for $m_{\eta'} = 1.6$ TeV, the diphoton rate is about $\sigma(pp \rightarrow \eta' \rightarrow \gamma\gamma) \simeq 0.98$ fb (for $\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma) = 6$ fb), so we expect a couple of events in the ATLAS data of 3.2 fb^{-1} , which is consistent with the “excess.” In this regard, another interesting possibility is that the G_H sector has an extra hidden quark that is singlet under the standard model gauge group. This is discussed in section 4.

The ρ_χ particle is pair produced by ordinary QCD interactions. The pair production cross section is expected to be of $O(10^{-1} - 10^{-9} \text{ fb})$ for the ρ_χ mass of 2 – 5 TeV at the 13 TeV LHC, although there may be deviations from this naive QCD estimate by a factor of a few due to a form factor. The produced ρ_χ decays into $\chi\psi$ or $\chi\varphi$, leading to two quasi-stable (or leptoquark-type) particles and 4 standard model gauge bosons per event. The decay of ρ_χ into $\chi\phi$ is forbidden by C -parity.

The particles ρ_φ and ρ_ϕ mix with the standard model $SU(2)_L$ and $U(1)_Y$ gauge bosons and couple to standard model quarks and leptons with couplings $\sim \sqrt{N}g_2^2/4\pi$ and $\sim \sqrt{N}g_1^2/4\pi$, respectively. The single production cross sections of ρ_φ and ρ_ϕ from quark initial states are of order 10 – 0.1 fb and 1 – 0.01 fb, respectively, for their masses of 2 – 5 TeV at the 13 TeV LHC. The coupling between ρ_ϕ and two gluons is absent due to C -parity. The produced ρ_φ decays mainly into $\varphi\varphi$, while the decay into $\varphi\phi$ is forbidden by C -parity. The ρ_ϕ decays into $\chi\chi$; the decays into $\psi\psi$, $\varphi\varphi$, and $\phi\phi$ are forbidden by C -parity.

The ω particle does not mix with the standard model gauge bosons. Furthermore, a coupling between ω and two gluons is forbidden by C -parity. Therefore, the dominant production of ω occurs through a coupling with three gluons

$$\mathcal{L} \approx \frac{\sqrt{N}g_3^3}{4\pi\Lambda^4}(\partial_\mu\omega_\nu)d^{abc}G^{a\nu\rho}G_{\rho\sigma}^bG^{c\sigma\mu}. \tag{2.26}$$

In the production process, the initial state is two gluons and the final state is a single ω and a gluon. The production cross section is expected to be of order 10 – 0.01 fb for the ω mass of 2 – 5 TeV at the 13 TeV LHC. The produced ω decays mainly into three hidden pions and to some extent into $\chi\chi$ (with the branching ratio of a few percent, suppressed by the size of the flavor $SU(5)$ breaking, $(m_D - m_L)^2/\Lambda^2$); the decays into $\psi\psi$, $\varphi\varphi$, and $\phi\phi$ are forbidden by C -parity and standard model gauge invariance.

3 Physics at higher energies

Some of the physics of hidden hadrons at the TeV scale are affected by theories at higher energies; for example, the lifetime of hidden pion χ is determined by the structure of the theory above Λ . Here we discuss particles that do not decay through G_H or standard model gauge dynamics, in particular χ and low-lying hidden baryons. We discuss low-energy operators necessarily to make these particles decay and study their phenomenological implications, including constraints from cosmology. We also discuss possible ultraviolet structures that lead to the required size of the coefficients of these operators.

3.1 Physics of hidden pion χ

The dynamics of G_H by itself leaves hidden pion χ absolutely stable. The χ particle, however, can decay into standard model particles through direct interactions between the G_H and standard model sectors. Here we study physics of χ decays, focusing on issues such

as bounds from cosmology and proton decay as well as implications for collider physics and theories at very high energies.⁴

At the level of the standard model fermion bilinears and hidden quark bilinears, the most general operators relevant for χ decays are given by

$$\mathcal{L} \sim \left\{ \begin{array}{c} \Psi_L \sigma_\mu \Psi_D^\dagger \\ \bar{\Psi}_D \sigma_\mu \bar{\Psi}_L^\dagger \end{array} \right\} \times \left\{ \begin{array}{c} q \sigma^\mu u^\dagger \\ e \sigma^\mu q^\dagger \\ d \sigma^\mu l^\dagger \end{array} \right\} + \text{h.c.}, \quad (3.1)$$

where $q(\square, \square, 1/6)$, $u(\bar{\square}, \mathbf{1}, -2/3)$, $d(\bar{\square}, \mathbf{1}, 1/3)$, $l(\mathbf{1}, \square, -1/2)$, and $e(\mathbf{1}, \mathbf{1}, 1)$ are the standard model (left-handed Weyl) fermions. Since the operators in the first bracket are matched on to $\partial_\mu \chi$, these operators give

$$\mathcal{L} = \alpha_1 (\partial_\mu \chi) (q \sigma^\mu u^\dagger) + \alpha_2 (\partial_\mu \chi) (e \sigma^\mu q^\dagger) + \alpha_3 (\partial_\mu \chi) (d \sigma^\mu l^\dagger) + \text{h.c.}, \quad (3.2)$$

at the scale Λ . Here, $\alpha_{1,2,3}$ are coefficients of order

$$\alpha_{1,2,3} \sim \frac{\sqrt{N}}{4\pi} \frac{\Lambda}{M_*^2}, \quad (3.3)$$

where M_* ($\gtrsim \Lambda$) is the scale at which the operators in eq. (3.1) are generated. We note that since $\Psi_L \sigma_\mu \Psi_D^\dagger$ and $\bar{\Psi}_D \sigma_\mu \bar{\Psi}_L^\dagger$ in eq. (3.1) correspond to conserved currents in the G_H sector, coefficients $\alpha_{1,2,3}$ are given by eq. (3.3) even if the G_H gauge theory is strongly coupled between M_* and Λ . With eqs. (3.2), (3.3), the decay width of χ is given by

$$\Gamma_\chi \approx \frac{1}{8\pi} \alpha_{1,2,3}^2 m_f^2 m_\chi \sim \frac{N}{(4\pi)^3} \frac{m_f^2 \Lambda^2}{M_*^4} m_\chi, \quad (3.4)$$

where m_f is the larger of the final state standard model fermion masses, arising from chirality suppression.

As we will see in section 3.2, cosmology requires the lifetime of χ to be smaller than $\approx O(10^{13}-10^{15} \text{ sec})$, assuming the standard thermal history below temperature of about a TeV. If the operators in eq. (3.2) are the only ones contributing to χ decays, then this gives

$$M_* \lesssim 10^{12}-10^{13} \text{ GeV} \times \sqrt{\frac{N}{5}} \left(\frac{m_f}{173 \text{ GeV}} \right)^{\frac{1}{2}} \left(\frac{\Lambda}{3.2 \text{ TeV} \sqrt{N/5}} \right)^{\frac{1}{2}} \frac{m_\chi}{1 \text{ TeV}}, \quad (3.5)$$

where we have used the last expression of eq. (3.4) and normalized m_f by the top quark mass. On the other hand, loops of hidden quarks involving eq. (3.1) induce standard model four-fermion operators suppressed by $\sim (4\pi M_*)^2$. If the coefficients of eq. (3.1) corresponding to α_1 and one of $\alpha_{2,3}$ are nonzero, this induces proton decay which is too rapid for the values of M_* satisfying eq. (3.5). This implies that if χ decays satisfy the cosmological bound due to eq. (3.2), then it must be that $\alpha_1 = 0$ or $\alpha_2 = \alpha_3 = 0$. In this case, M_* may be small enough that χ decays within the detector (even promptly).

⁴Physics of a composite pseudo Nambu-Goldstone boson that has the same standard model gauge quantum numbers as χ was discussed in [45].

One might think that the upper bound on M_* in eq. (3.5) implies that unless early universe cosmology is exotic, there must be new physics directly connecting the G_H and standard model sectors well below the conventional unification scale, $\sim 10^{14}\text{--}10^{17}$ GeV. This is, however, not necessarily true. If neither $\Psi_{D,L}$ nor $\bar{\Psi}_{D,L}$ is unified in a single multiplet at the high energy scale (as in the case of higher dimensional grand unified theories [46]), we can consider operators leading to χ decays involving standard model operators of dimension 4:

$$\mathcal{L} \sim \left\{ \begin{array}{l} \bar{\Psi}_D \Psi_L \\ \Psi_D^\dagger \bar{\Psi}_L^\dagger \end{array} \right\} \times \left\{ \begin{array}{ll} qqh^\dagger, & ueh^\dagger \\ deh, & q^\dagger l^\dagger h^\dagger \\ u^\dagger u^\dagger h, & u^\dagger d^\dagger h^\dagger \end{array} \right\} + \text{h.c.}, \quad (3.6)$$

with $O(1)$ coefficients at the unification scale M_* . (If $\Psi_{D,L}$ or $\bar{\Psi}_{D,L}$ is unified, operators of this form with $O(1)$ coefficients lead to too large mass terms for the hidden quarks at the radiative level.) These operators lead to operators at the scale Λ

$$\mathcal{L} = \beta_1 \chi qqh^\dagger + \beta_2 \chi ueh^\dagger + \beta_3 \chi deh + \beta_4 \chi q^\dagger l^\dagger h^\dagger + \beta_5 \chi u^\dagger u^\dagger h + \beta_6 \chi u^\dagger d^\dagger h^\dagger + \text{h.c.}, \quad (3.7)$$

where $h(\mathbf{1}, \square, -1/2)$ is the standard model Higgs field, and $\beta_{1,\dots,6}$ are coefficients. An important point is that since the operators in the first bracket in eq. (3.6) are scalar, their dimensions can receive large quantum corrections if the G_H theory is strongly coupled below M_* . Suppose that the G_H theory is in a conformal phase between M_* and M_I ($\gtrsim \Lambda$) because of the existence of extra matter of mass M_I charged under G_H . In this case, the size of the coefficients in eq. (3.7) is estimated as

$$\beta_{1,\dots,6} \sim \frac{\sqrt{N}}{4\pi} \frac{\Lambda^2}{M_*^3} \left(\frac{M_*}{M_I} \right)^{3-\Delta}, \quad (3.8)$$

where Δ (> 1) is the operator dimension in the conformal phase, which we take to be the same for $\bar{\Psi}_D \Psi_L$ and $\Psi_D^\dagger \bar{\Psi}_L^\dagger$ for simplicity. (They can in general be different depending on the structure of the theory in the energy interval between M_* and M_I .) It is therefore possible that the cosmological bound on χ decays $\tau_\chi \lesssim O(10^{13}\text{--}10^{15}$ sec), which translates into $\beta_{1,\dots,6}^{-1} \lesssim 10^{21}\text{--}10^{22}$ GeV, is satisfied even if M_* is at the unification scale. For example, if M_I is within an order of magnitude of Λ , the bound can be satisfied for $\Delta \lesssim 1.2 - 1.7$, with the precise value depending on M_* , M_I , and so on.

The large size of the anomalous dimension leading to small Δ requires that G_H is strongly coupled in the conformal phase. On the other hand, if this phase is still the G_H gauge theory, operators of the form

$$(\bar{\Psi}\Psi)(\bar{\Psi}\Psi)^\dagger, \quad (3.9)$$

must have dimensions larger than or equal to 4. Since these dimensions are given by 2Δ in the large N limit (at least for the ones directly related to $\bar{\Psi}_D \Psi_L$ or $\bar{\Psi}_L \Psi_D$ but also for others if the conformal phase respects the $SU(5)$ flavor symmetry), an $O(1/N)$ correction must play the role for making Δ smaller than 2. This seems to imply that it is difficult to

have M_* in the large side, $M_* \approx 10^{17}$ GeV, which requires $\Delta \lesssim 1.2$. Furthermore, strong coupling is also expected to give large anomalous dimensions for operators of the form

$$(\bar{\Psi}\Psi)^2. \tag{3.10}$$

The dimensions of these operators must also be larger than 4 unless their coefficients are strongly suppressed at M_* , since otherwise the theory flows into a different phase above the scale Λ , so that our low energy theory is not the one discussed in section 2. Avoiding the cosmological bound on χ using only operators in eq. (3.6) generated at a unification scale requires these conditions to be satisfied.

3.2 Cosmology

Assuming that the thermal history of the universe is standard below the temperature of about a TeV, the relic abundance of χ is given as follows. Below the temperature $T \simeq m_\chi/30 \simeq 30 - 40$ GeV, annihilation of χ into gluons becomes ineffective and the χ abundance freezes out. If there were no subsequent annihilation of χ , this would lead to the present-day χ energy density of $\Omega_\chi \approx O(0.01)$. However, we expect a period of further annihilation at $T \simeq \Lambda_{\text{QCD}} \approx O(100 \text{ MeV})$ when nonperturbative QCD effects become important. At $T \simeq \Lambda_{\text{QCD}}$, χ particles hadronize by picking up light quarks, which leads to enhanced annihilation of χ . While the details of this late-time annihilation process are not fully understood, we may estimate, based on earlier work [47, 48], that the resulting χ energy density is of order

$$10^{-8} \lesssim \Omega_\chi \lesssim 10^{-6}. \tag{3.11}$$

With this estimate, the possibility of absolutely stable χ is excluded, under reasonable assumptions, by heavy isotope searches; see section VC of ref. [45] (in which a particle called a xyon corresponds to our χ particle here).

For unstable χ , bounds on the lifetime are given by the requirement that cosmological decay of χ does neither spoil the success of the big bang nucleosynthesis, generate a detectable level of (μ or y) distortions of the cosmic microwave background, or lead to an excessive amount of the diffuse gamma-ray background. These constraints, as compiled in [49] with the update to include [50], are plotted in figure 4 in the τ_χ - m_χ - Y_χ plane, where τ_χ , m_χ , and $Y_\chi = \rho_\chi/s$ are the lifetime, mass, and entropy yield of χ . The estimate of the χ abundance after the QCD era (corresponding to eq. (3.11) if χ were absolutely stable) is indicated by the horizontal band. We thus find that with this estimate, the bound on the χ lifetime comes from observations of the diffuse gamma-ray background and is given by

$$\tau_\chi \lesssim 10^{13}\text{--}10^{15} \text{ sec}. \tag{3.12}$$

Implications of this bound were discussed in section 3.1.

In addition to the χ particle, the dynamics of G_H leaves low-lying hidden baryons stable. Among all the hidden baryons, we expect the lightest ones to be those in which the constituents form the smallest spin, i.e. 0 or 1/2, and have a vanishing orbital angular momentum [51].⁵ These states have masses $\sim N\Lambda$ and lighter than the other hidden

⁵Here we assume that effects involving both orbital and spin angular momenta do not invalidate this expectation, which is indeed the case in QCD.

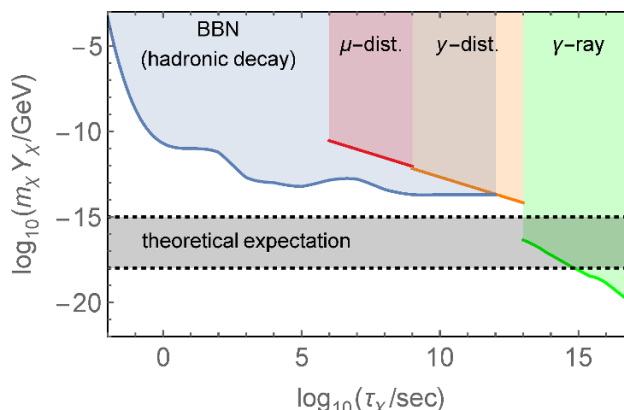


Figure 4. Cosmological constraints on the lifetime, τ_χ , and the product of the mass, m_χ , and entropy yield, Y_χ of the χ particle. The shaded regions are excluded by the analyses of big bang nucleosynthesis (BBN), μ and y distortions of the cosmic microwave background, and the diffuse gamma-ray background. The theoretical estimate of $m_\chi Y_\chi$ is indicated by the horizontal band.

baryons by $\sim \Lambda$, so all the heavier hidden baryons decay into these low-lying hidden baryons by emitting hidden pions and standard model gauge bosons. On the other hand, decays of the low-lying hidden baryons require interactions beyond G_H and standard model gauge interactions. In the limit of $m_D = m_L$ and vanishing standard model gauge couplings, the low-lying hidden baryons form an $SU(5)$ multiplet and are degenerate in mass. Mass splittings among them, therefore, are of order $|m_D - m_L| \approx O(100 \text{ GeV})$ or $g_{1,2,3}^2 N \Lambda / 16 \pi^2 \approx O(10 - 100 \text{ GeV})$, so that decays among low-lying hidden baryons cannot occur by emitting on-shell χ particles. This implies that the dynamics of G_H itself leaves all the low-lying hidden baryons stable.

The cosmological fate of the low-lying hidden baryons depends on the strength of operators violating hidden baryon number and operators responsible for χ decays. The former operators lead to decays of hidden baryons either directly to standard model particles or to χ and standard model particles through other (off-shell) hidden baryons. If the timescale for these decays, τ_B , is shorter than the χ lifetime, τ_χ , then all the low-lying hidden baryons decay in this manner. On the other hand, if this timescale is longer than τ_χ , the heavier components of the low-lying hidden baryons decay first into the lightest hidden baryon (and standard model particles) in timescale τ_χ ; then the lightest hidden baryon decays to standard model particles (and often χ) in a longer timescale of τ_B .

Cosmology of hidden baryons is N dependent, since the spectrum of the low-lying hidden baryons as well as operators responsible for their decays strongly depend on N . There are, however, some statements one can make regardless of the value of N . Suppose $m_D > m_L$. (The other case is discussed later.) In this case, among the low-lying hidden baryons the one composed only of Ψ_L is the lightest. Then, for even and odd N the lightest hidden baryon has standard model gauge quantum numbers of $(\mathbf{1}, \mathbf{1})_{-N/2}$ and $(\mathbf{1}, \square)_{-N/2}$, respectively, and hence is electrically charged. The thermal abundance of the lightest hidden baryon is determined by its annihilation into hidden pions and given by $n_B/s \sim 10^{-16} \times (m_B/10 \text{ TeV})$. An electrically charged particle with such an abundance

is excluded by searches for charged massive stable particles [52, 53]. Thus, the lightest hidden baryon must be unstable in this case.

We now discuss physics of hidden baryon decays in more detail. For illustrative purposes, here we consider the cases of $N = 3, 4, 5$ for $m_D > m_L$. The cases of higher N can be analyzed analogously.

$N = 3$: the lightest hidden baryon, which is LLL in obvious notation, can decay via interactions of the form

$$\mathcal{L} \sim u\bar{\Psi}_D\bar{\Psi}_L\bar{\Psi}_L + q\bar{\Psi}_D\bar{\Psi}_D\bar{\Psi}_L + e\bar{\Psi}_D\bar{\Psi}_D\bar{\Psi}_D + \text{h.c.}, \quad (3.13)$$

which induce mixings between the DLL , DDL , and DDD hidden baryons and the right-handed up-type quarks, left-handed quarks, and right-handed leptons, respectively:

$$\mathcal{L} = \gamma_1 u\bar{B}_{DLL} + \gamma_2 q\bar{B}_{DDL} + \gamma_3 e\bar{B}_{DDD} + \text{h.c.}, \quad \gamma_{1,2,3} \sim \frac{\sqrt{N}}{(4\pi)^2} \frac{\Lambda^3}{M_*^2} \left(\frac{M_*}{M_I}\right)^{9/2-\Delta_B}. \quad (3.14)$$

Here, we have included possible enhancement of the coefficients by conformal dynamics between the scales M_* and M_I , with Δ_B ($> 3/2$) representing the dimension of hidden baryonic operators $\bar{\Psi}\bar{\Psi}\bar{\Psi}$ in the conformal phase.⁶ The lightest hidden baryon LLL then decays into χ and standard model fermions with the decay rate

$$\Gamma_{LLL} \sim \frac{4\pi\gamma_{1,2,3}^2}{\Lambda}. \quad (3.15)$$

Since this state decays hadronically with the abundance of $\rho_B/s \sim 10^{-12}$ GeV \times $(m_B/10 \text{ TeV})^2$, preserving the success of the big bang nucleosynthesis requires the lifetime to be shorter than $O(100 \text{ sec})$ [54], which translates into $\gamma_{1,2,3} \gtrsim 10^{-12}$ GeV. For $M_* \sim 10^{14}\text{--}10^{17}$ GeV and $M_I \sim \Lambda$, this requires $\Delta_B \lesssim 3.4 - 3.7$.

The decay of the lightest hidden baryon produces χ . If $\tau_{LLL} \equiv \Gamma_{LLL}^{-1} < 10^{-4}$ sec, the produced χ is subject to QCD enhanced annihilation afterward, so that the analysis in the previous subsection persists. On the other hand, if $\tau_{LLL} > 10^{-4}$ sec, the abundance of χ is determined by its annihilation just after the decay of the lightest hidden baryon, which is roughly given by

$$\frac{\rho_\chi}{s} \sim m_\chi \frac{1}{(\Gamma_{LLL} M_{\text{Pl}})^{3/2}} \frac{\Gamma_{LLL}}{\sigma_{\chi\bar{\chi}}} \simeq 10^{-13} \text{ GeV} \frac{m_\chi}{1 \text{ TeV}} \left(\frac{\tau_{LLL}}{100 \text{ sec}}\right)^{1/2} \frac{1 \text{ fm}^2}{\sigma_{\chi\bar{\chi}}}, \quad (3.16)$$

where M_{Pl} is the reduced Planck scale. The bounds on τ_χ in this case can be read off from figure 4.

We note that the hidden baryon number violating operators discussed here also induce decays of χ as they violate D and L numbers. Through the mixing between hidden baryons and standard model fermions, χ decays into uq^\dagger or qe^\dagger with the decay rate

$$\Gamma_\chi \sim 4\pi N m_\chi \left(\frac{\gamma_{1,2,3}}{N\Lambda}\right)^4. \quad (3.17)$$

⁶In our analysis here, we assume that the dimensions of different hidden baryonic operators are (approximately) the same. In particular, we assume that the differences of the dimensions are small enough that operators responsible for hidden baryon decays are given by those in which the sum of the canonical dimensions of the standard model fields is the smallest.

If $\gamma_{1,2,3} \gtrsim 10^{-8} - 10^{-7}$ GeV, the lifetime of χ can be shorter than $O(10^{13} - 10^{15}$ sec) without having the operators discussed in the previous section. For $M_* \sim 10^{14} - 10^{17}$ GeV and $M_I \sim \Lambda$, this translates into $\Delta_B \lesssim 3.1 - 3.3$. For such large $\gamma_{1,2,3}$, hidden baryons are short-lived and do not cause any cosmological problems.

$N = 4$: the lightest hidden baryon, $LLLL$, can decay via

$$\mathcal{L} \sim h \Psi_D \Psi_D \Psi_D \Psi_L + \text{h.c.} \quad (3.18)$$

This induces a mixing between the $DDDL$ hidden baryon and the standard model Higgs field

$$\mathcal{L} = \delta h B_{DDDL} + \text{h.c.}, \quad \delta \sim \frac{\sqrt{N}}{(4\pi)^3} \frac{\Lambda^5}{M_*^3} \left(\frac{M_*}{M_I} \right)^{6-\Delta_B}, \quad (3.19)$$

where $\Delta_B (> 1)$ is the dimension of the operator $\Psi_D \Psi_D \Psi_D \Psi_L$ in the conformal phase. The lightest hidden baryon then decays into three χ and a standard model Higgs or gauge boson, with the decay rate

$$\Gamma_{LLLL} \sim 4\pi \frac{\delta^2}{N^2 \Lambda^3}. \quad (3.20)$$

As the standard model Higgs and gauge bosons decay hadronically, the lifetime of the lightest hidden baryon must be shorter than $O(100$ sec), which requires $\delta \gtrsim 10^{-8}$ GeV². For $M_* \sim 10^{14} - 10^{17}$ GeV and $M_I \sim \Lambda$, this translates into $\Delta_B \lesssim 3.8 - 4.0$. Cosmology of χ produced by the lightest hidden baryon decay is as in the case of $N = 3$.

$N = 5$: the lightest hidden baryon, $LLLLL$, can decay via interactions of the form

$$\mathcal{L} \sim h^\dagger d \Psi_D \Psi_D \Psi_L \Psi_L \Psi_L + h^\dagger l \Psi_D \Psi_D \Psi_D \Psi_L \Psi_L + \text{h.c.} \quad (3.21)$$

These operators induce couplings between the $DDLLL$ and $DDDLL$ hidden baryons with standard model particles

$$\mathcal{L} = \epsilon_1 h^\dagger d B_{DDLLL} + \epsilon_2 h^\dagger l B_{DDDLL}, \quad \epsilon_{1,2} \sim \frac{\sqrt{N}}{(4\pi)^4} \frac{\Lambda^6}{M_*^6} \left(\frac{M_*}{M_I} \right)^{15/2-\Delta_B}, \quad (3.22)$$

where $\Delta_B (> 3/2)$ is the dimension of hidden baryonic operators $\Psi \Psi \Psi \Psi \Psi$ in the conformal phase. The lightest hidden baryon decays into a standard model Higgs or gauge boson, a down quark or lepton doublet, and multiple χ , with the decay rate

$$\Gamma_{LLLLL} \sim \frac{N}{4\pi} \epsilon_{1,2}^2 \Lambda. \quad (3.23)$$

As the decay is hadronic, the lifetime of the $LLLLL$ hidden baryon must be shorter than $O(100$ sec), which requires $\epsilon_{1,2} \gtrsim 10^{-15}$. For $M_* \sim 10^{14} - 10^{17}$ GeV and $M_I \sim \Lambda$, this translates into $\Delta_B \lesssim 2.3 - 2.5$. Cosmology of χ produced by the decay can be analyzed as in the case of $N = 3$.

We now discuss the case with $m_D < m_L$. In this case, depending on N and the precise values of $m_{D,L}$, the lightest hidden baryon may carry standard model color. Since a colored hidden baryon is subject to late-time annihilation around the QCD phase transition era,

the bound on its lifetime is weak, $\tau \lesssim O(10^{13}\text{--}10^{15} \text{ sec})$. This may allow for the coefficients of the hidden baryon number violating operators to be much smaller than the case discussed above. If χ decay is prompt, this is indeed the case because then all the heavier low-lying hidden baryons decay into the lightest hidden baryon before the QCD annihilation era. On the other hand, if χ is long-lived, decays of the low-lying hidden baryons are all controlled by the hidden baryon number violating operators. The bounds on the coefficients are then (essentially) the same as before, since they are determined by the decays of non-colored hidden baryons.

As seen here, cosmology of hidden baryons is controlled by the lowest dimensional hidden baryon number violating operator. Since its dimension depends on the whole hidden quark content, the existence of an extra hidden quark could alter the situation. In particular, if there is an extra hidden quark charged under G_H but singlet under G_{SM} , then the decays of hidden baryons can be faster. This case will be discussed in section 4.

3.3 Possible ultraviolet theories

We have discussed physics of “would-be” stable particles: χ and low-lying hidden baryons. Assuming the standard thermal history of the universe below the TeV scale, we have found that these particles must decay fast enough via non-renormalizable interactions. The required strength of these interactions suggests the existence of ultraviolet physics beyond the minimal G_H and standard model sectors not too far above the dynamical scale, Λ .⁷

One possibility is that physics responsible for the χ and hidden baryon decays is indeed at a scale M_* which is within a few orders of magnitudes of Λ . Here we consider an alternative possibility that physics leading to these decays is at very high energies, e.g. at a unification scale $M_* \approx O(10^{14}\text{--}10^{16} \text{ GeV})$. Here, we consider a few examples that this can be the case.

- *Conformal dynamics:* as we have discussed before, even if the relevant non-renormalizable interactions are generated far above Λ , such as at the unification scale, conformal dynamics of G_H may enhance these operators and induce sufficiently fast decays of the would-be stable particles. Having conformal dynamics requires a sufficient number of (vectorlike) particles charged under G_H in addition to $\Psi_{D,L}$ and $\bar{\Psi}_{D,L}$, which we may assume to have masses larger than the dynamical scale by a factor of $O(1 - 10)$. In this case, we can consider that G_H is in a conformal phase above the mass scale of these particles but deviates from it below this mass scale and finally confines at Λ . With such dynamics, we can understand the proximity of the dynamical scale of G_H , $\Lambda \sim \text{TeV}$, and the masses of hidden quarks $m \sim 0.1 \text{ TeV}$, if the masses $\Psi_{D,L}$, $\bar{\Psi}_{D,L}$ and additional particles originate from a common source and the conformal phase of G_H is strongly coupled. The existence of a particle charged under G_H and singlet under G_{SM} may also help satisfying the constraints from cosmology; see section 4. We note that (some of) the additional particles added here

⁷An alternative possibility is that the thermal history of the early universe is non-standard. For detailed analyses of the relic abundance of quasi-stable particles in the case that the reheating temperature is very low, see e.g. [55].

may be charged under G_{SM} , which does not destroy gauge coupling unification if they form complete $SU(5)$ multiplets.

- *Supersymmetry*: if superpartners of the standard model particles and hidden quarks are near the TeV scale, then the decay rates of would-be stable particles can be larger than those in the non-supersymmetric model. (With the superpartners near the TeV scale, $N \leq 4$ is required in order for the standard model gauge couplings not to blow up below the unification scale.) Specifically, hidden baryon number can be more easily broken. For $N = 3$, for example, there are dimension-five superpotential operators

$$W \sim \frac{1}{M_*} (u\bar{\Psi}_D\bar{\Psi}_L\bar{\Psi}_L + q\bar{\Psi}_D\bar{\Psi}_D\bar{\Psi}_L + e\bar{\Psi}_D\bar{\Psi}_D\bar{\Psi}_D), \quad (3.24)$$

which mix standard model particles with hidden baryons. Even if M_* is around the unification scale, hidden baryons decay with lifetimes shorter than $O(100 \text{ sec})$ if superpartner masses are around $O(1 - 10 \text{ TeV})$. Introduction of superparticles at the TeV scale also improves gauge coupling unification.

- *Superconformal dynamics*: it is possible that both conformal dynamics and supersymmetry are at play above some scale not far from Λ . This scenario is very interesting. Conformal dynamics as well as exchange of superpartners may enhance the decay rates of would-be stable particles, and the deviation of G_H from the conformal window may be explained by the decoupling of hidden squarks and hidden gauginos at $O(1 - 10 \text{ TeV})$. Moreover, the anomalous dimensions of some of the operators can be calculated due to supersymmetry.

The ultraviolet structure of this class of theories is tightly constrained if the G_H theory is already in a conformal phase at the unification scale M_* . Let us first assume that all the superpartners are at $\tilde{m} \approx O(\text{TeV})$. In this case, $N \geq 4$ leads to a Landau pole for the standard model gauge couplings below M_* . This is because in the energy interval between \tilde{m} and M_* , the G_H gauge theory is on a nontrivial fixed point, so that the effect of G_H gauge interactions enhances the contribution of the hidden quarks to the running of the standard model gauge couplings (at the level of two loops in the standard model gauge couplings). Even for $N = 3$, the number of flavors of the G_H gauge theory is bounded from below, since the G_H gauge coupling at the fixed point is larger for a smaller number of flavors, making the contribution from the hidden quarks larger. We find that the number of flavors must be 7 or 8 so that G_H is in the conformal window while at the same time the standard model gauge couplings do not hit a Landau pole below the unification scale. Assuming that additional particles always appear in the form of complete $SU(5)$ multiplets, this predicts the existence of particles that are charged under G_H but singlet under G_{SM} . As we will see in section 4, the existence of such particles helps to evade constraints from cosmology. For heavier superparticle masses, the possible choice of N and the number of flavors, F , increases; for example, for $\tilde{m} \approx O(10 \text{ TeV})$, $(N, F) = (3, 6)$, $(4, 10)$, and $(4, 11)$ are also allowed.

	$G_H = \text{SU}(N)$	$\text{SU}(3)_C$	$\text{SU}(2)_L$	$\text{U}(1)_Y$
Ψ_D	\square	$\bar{\square}$	$\mathbf{1}$	$1/3$
Ψ_L	\square	$\mathbf{1}$	\square	$-1/2$
Ψ_N	\square	$\mathbf{1}$	$\mathbf{1}$	0
$\bar{\Psi}_D$	$\bar{\square}$	\square	$\mathbf{1}$	$-1/3$
$\bar{\Psi}_L$	$\bar{\square}$	$\mathbf{1}$	\square	$1/2$
$\bar{\Psi}_N$	$\bar{\square}$	$\mathbf{1}$	$\mathbf{1}$	0

Table 2. Charge assignment of the extended model. $\Psi_{D,L,N}$ and $\bar{\Psi}_{D,L,N}$ are left-handed Weyl spinors.

4 An extra hidden light quark

In this section, we discuss a singlet extension of the model described in section 2: we add an extra vectorlike hidden quark, Ψ_N and $\bar{\Psi}_N$, which is in the fundamental representation of G_H but is singlet under the standard model gauge group. As mentioned in section 3.2, this makes it easier to avoid constraints from cosmology. It also leads to two diphoton resonances in the spectrum of hidden pions, which has interesting phenomenological implications (one of which was mentioned in footnote 3). The full matter content of the model is summarized in table 2. This charge assignment was also considered in [16]. We here take the masses of the hidden quarks

$$\mathcal{L} = -m_D \Psi_D \bar{\Psi}_D - m_L \Psi_L \bar{\Psi}_L - m_N \Psi_N \bar{\Psi}_N + \text{h.c.}, \tag{4.1}$$

to be real and positive without loss of generality. We assume $m_{D,L,N} \lesssim \Lambda$.

The spectrum below Λ consists of hidden pions

$$\begin{aligned} \psi(\mathbf{Adj}, \mathbf{1}, 0), \quad \chi\left(\square, \square, -\frac{5}{6}\right), \quad \varphi(\mathbf{1}, \mathbf{Adj}, 0), \quad \phi(\mathbf{1}, \mathbf{1}, 0), \\ \xi\left(\square, \mathbf{1}, -\frac{1}{3}\right), \quad \lambda\left(\mathbf{1}, \square, \frac{1}{2}\right), \quad \eta(\mathbf{1}, \mathbf{1}, 0), \end{aligned} \tag{4.2}$$

where χ , ξ , and λ are complex while the others are real. Note that there are two hidden pions which are singlet under the standard model gauge group, G_{SM} —one is charged under $\text{SU}(5) \supset G_{\text{SM}}$ while the other is singlet, which we refer to as ϕ and η , respectively. The masses of the hidden pions are given by

$$m_\psi^2 = 2m_D \frac{c}{f^2} + 3\Delta_C, \tag{4.3}$$

$$m_\chi^2 = (m_D + m_L) \frac{c}{f^2} + \frac{4}{3}\Delta_C + \frac{3}{4}\Delta_L + \frac{5}{12}\Delta_Y, \tag{4.4}$$

$$m_\phi^2 = 2m_L \frac{c}{f^2} + 2\Delta_L, \tag{4.5}$$

$$m_\xi^2 = (m_D + m_N) \frac{c}{f^2} + \frac{4}{3}\Delta_C + \frac{1}{15}\Delta_Y, \tag{4.6}$$

$$m_\lambda^2 = (m_L + m_N) \frac{c}{f^2} + \frac{3}{4} \Delta_L + \frac{3}{20} \Delta_Y, \quad (4.7)$$

$$\begin{pmatrix} m_\phi^2 & m_{\phi\eta}^2 \\ m_{\eta\phi}^2 & m_\eta^2 \end{pmatrix} = \begin{pmatrix} \frac{2}{5}(2m_D + 3m_L) & \frac{2}{5}(m_D - m_L) \\ \frac{2}{5}(m_D - m_L) & \frac{1}{15}(3m_D + 2m_L + 25m_N) \end{pmatrix} \frac{c}{f^2}, \quad (4.8)$$

where c and f are the hidden quark bilinear condensates and the decay constant, respectively.

The mixing between ϕ and η vanishes for $m_D = m_L$ due to the SU(5) symmetry. Expect for this special case, the mass eigenstates ϕ_+ and ϕ_- are determined by eq. (4.8)

$$\phi_+ = \eta \cos \theta + \phi \sin \theta, \quad \phi_- = -\eta \sin \theta + \phi \cos \theta. \quad (4.9)$$

Here, the mixing angle θ and the mass eigenvalues m_+ and m_- ($m_+^2 > m_-^2$) are related with $m_{D,L,N}$ as

$$m_D = \frac{f^2}{c} \frac{m_-^2 - 3(m_+^2 - m_-^2) \tan \theta + m_+^2 \tan^2 \theta}{2(1 + \tan^2 \theta)}, \quad (4.10)$$

$$m_L = \frac{f^2}{c} \frac{m_-^2 + 2(m_+^2 - m_-^2) \tan \theta + m_+^2 \tan^2 \theta}{2(1 + \tan^2 \theta)}, \quad (4.11)$$

$$m_N = \frac{f^2}{c} \frac{6m_+^2 - m_-^2 + (m_+^2 - m_-^2) \tan \theta + (6m_-^2 - m_+^2) \tan^2 \theta}{10(1 + \tan^2 \theta)}. \quad (4.12)$$

The dimension-five couplings of the hidden pions with the standard model gauge fields are determined by chiral anomalies. The couplings of ψ , φ , and ϕ are given by eq. (2.10), while those of η are given by eq. (2.22) with the replacement $\eta' \rightarrow \eta/\sqrt{6}$. The couplings of the mass eigenstates ϕ_\pm can be read off from these expressions and the mixing in eq. (4.9).

4.1 Diphoton (diboson) signals and other phenomenology

A distinct feature of this model is that there are two standard model singlet hidden pions, ϕ_+ and ϕ_- , which are produced via gluon fusion and decay into a pair of standard model gauge bosons, including a diphoton. If kinematically allowed, they may also decay into three hidden pions; with parity violation by $\theta_H \neq 0$, they also decay into two hidden pions. Here we assume that the decay channels into hidden pions are suppressed kinematically and/or by $\theta_H \simeq 0$. There are several interesting possibilities to consider in terms of the phenomenology of these particles:

- *1.6 TeV diphoton excess.* We may identify the two singlets ϕ_- and ϕ_+ as the origins of, respectively, the 750 GeV excess and the slight “excess” at $\simeq 1.6$ TeV seen in the ATLAS diphoton data [1, 2]. In the left panel of figure 5, we show the values of the hidden quark masses $m_{D,L,N}$ needed to reproduce $m_- = 750$ GeV and $m_+ = 1.6$ TeV as a function of θ . Here, the decay constant f has been determined so that $\sigma(pp \rightarrow \phi_- \rightarrow \gamma\gamma) = 6$ fb at $\sqrt{s} = 13$ TeV, as plotted in the right panel. In the parameter region $\theta/\pi \in [0, 0.03) \cup (0.37, 0.70) \cup (0.95, 1)$, $m_- = 750$ GeV and $m_+ = 1.6$ TeV can be obtained by an appropriate choice of $m_{D,L,N}$, but in the other regions — which are shaded — no choice of $m_{D,L,N}$ can lead to the required ϕ_\pm masses. Near the

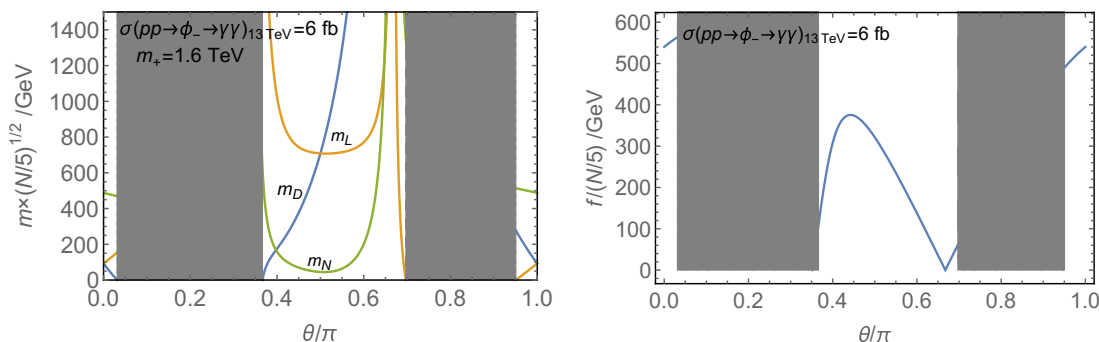


Figure 5. The hidden quark masses, m_D , m_L and m_N , reproducing $m_- = 750$ GeV and $m_+ = 1.6$ TeV as a function of θ (left). Here, the value of the decay constant f is determined so that $\sigma(pp \rightarrow \phi_- \rightarrow \gamma\gamma) = 6$ fb is obtained at $\sqrt{s} = 13$ TeV (right). In the gray-shaded regions of θ , no choice of $m_{D,L,N}$ may reproduce the required ϕ_{\pm} masses.

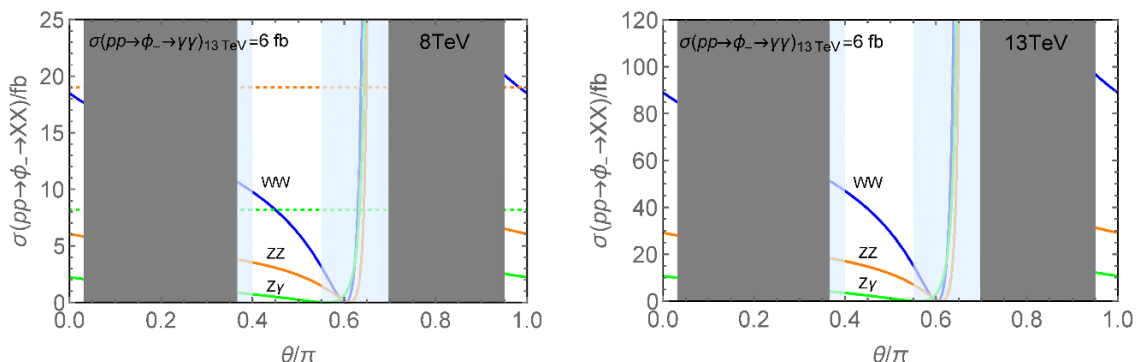


Figure 6. The production cross section of ϕ_- times branching ratios into two electroweak gauge bosons at $\sqrt{s} = 8$ TeV (left) and 13 TeV (right). The gray-shaded regions are unphysical. In the blue-shaded regions, chiral perturbation theory cannot be fully trusted, and the plots represent only qualitative estimations.

edge of the allowed region $(0.37, 0.70)$, the decay constant f is required to be small. This is because at $\theta/\pi \sim 0.35$ the dimension-five coupling between ϕ_- and gluons vanishes, while at $\theta/\pi \sim 0.65$ that between ϕ_- and photons vanishes, so formally $f \rightarrow 0$ is required to obtain $\sigma(pp \rightarrow \phi_- \rightarrow \gamma\gamma) = 6$ fb at these values of θ . Note that for $m_{D,L,N} \gtrsim \Lambda \simeq 4\pi f/\sqrt{N}$, the results obtained here using chiral perturbation theory expressions, eq. (4.3)–(4.8), cannot be fully trusted, although we may still regard them as giving qualitatively correct estimates.

In figure 6, we present predictions for the production cross section of ϕ_- times the branching ratios into two electroweak gauge bosons at $\sqrt{s} = 8$ TeV (left) and 13 TeV (right). In the blue-shaded regions, chiral perturbation theory gives only qualitative estimates because of $m_{D,L,N} \gtrsim \Lambda$. In the left panel, we also show the upper bound on the cross section for each mode [5, 7, 33] by the dotted line using the same color as the corresponding prediction. For any $\theta/\pi \in [0, 0.03) \cup (0.4, 0.55) \cup (0.95, 1)$, in which chiral perturbation theory can be trusted, the bounds are all satisfied. Predictions

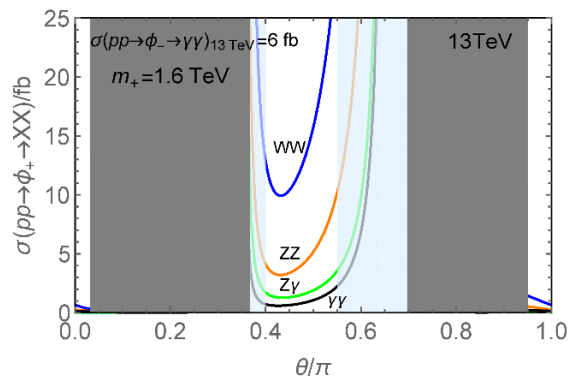


Figure 7. The production cross section of ϕ_+ of mass $m_+ = 1.6$ TeV times the branching ratios into two electroweak gauge bosons at $\sqrt{s} = 13$ TeV.

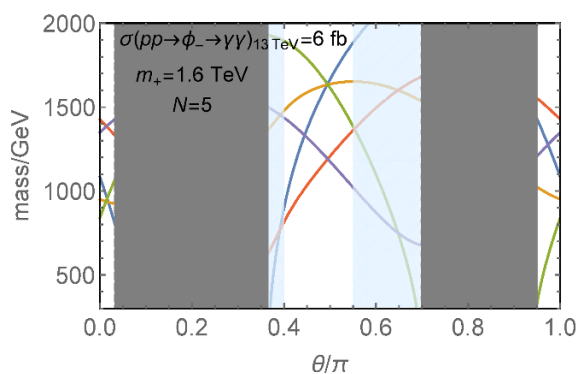


Figure 8. The masses of hidden pions ψ (blue), χ (orange), φ (green), ξ (red) and λ (purple) for $m_- = 750$ GeV and $m_+ = 1.6$ TeV as a function of θ . Here, we have chosen $N = 5$.

for the production cross section of the other neutral hidden pion, ϕ_+ , times branching ratios into two electroweak gauge bosons at $\sqrt{s} = 13$ TeV are presented in figure 7. For $\theta \sim \pi/2$, the production cross section of diphoton is $O(\text{fb})$, which is consistent with the “excess” at $\simeq 1.6$ TeV. We also show the prediction for the masses of the hidden pions in figure 8. For $\theta \sim \pi/2$, the mass of the χ particle can reach ~ 1.5 TeV. This helps to evade the experimental bound on this particle (see section 2.2).

- *(Apparent) wide width of the 750 GeV excess.* Alternatively, we may consider that both ϕ_- and ϕ_+ have masses around 750 GeV with a small mass difference of 10s of GeV. In this case, the two resonances are observed as an apparent wide resonance [7, 56], which is mildly preferred by the ATLAS diphoton data. Such mass degeneracy occurs if $m_D \simeq m_L \simeq m_N$. The required value of the decay constant f is larger than that in eq. (2.13):

$$f \simeq 720 \text{ GeV} \frac{N}{5} \sqrt{\frac{6 \text{ fb}}{\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma)}}, \quad (4.13)$$

because the two resonances contribute to the diphoton rate. With eqs. (4.3)–(4.8), (4.13) and $m_D \simeq m_L \simeq m_N$, the masses of the hidden pions are determined for

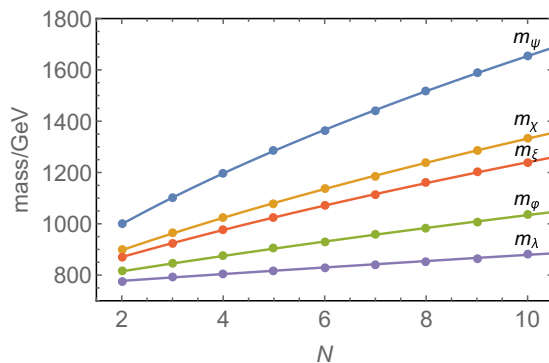


Figure 9. The masses of hidden pions ψ , χ , φ , ξ , and λ for $m_- \simeq m_+ \simeq 750$ GeV for given values of N .

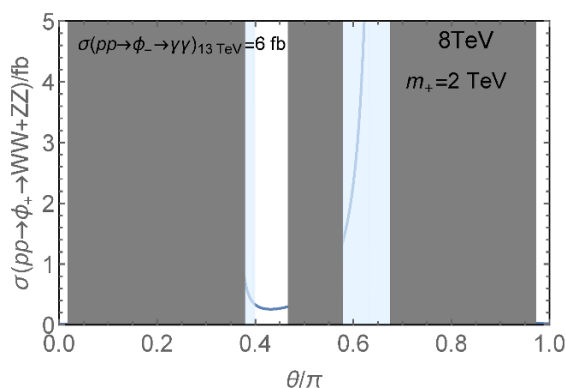


Figure 10. The production cross section of ϕ_+ of mass $m_+ = 2$ TeV times the sum of the branching ratios into WW and ZZ at $\sqrt{s} = 8$ TeV.

given N , which are shown in figure 9. Since f is larger, the hidden pion masses are larger than in the model without the singlet hidden quark. In particular, m_χ easily exceeds 1 TeV, helping to evade the constraint.

- *2 TeV diboson excess.* Yet another possibility is that ϕ_- and ϕ_+ have masses $\simeq 750$ GeV and $\simeq 2$ TeV, with the latter responsible for the diboson excess reported by the ATLAS collaboration [21] through the WW and ZZ decay modes [19, 20]. In figure 10, we show the production cross section of WW and ZZ through ϕ_+ . The meaning of the shades is the same as in figure 6. We find that the cross section of a few fb, which is needed to explain the excess, may be achieved for $\theta/\pi \sim 0.6$, although chiral perturbation theory can give only qualitative estimates in this region.

We note that in addition to ϕ_- and ϕ_+ , the model also has η' associated with the anomalous U(1) axial flavor symmetry. If its mass is somewhat lower than the naive expectation and if the G_H sector respects CP , then we have three scalar resonances decaying into two standard model gauge bosons in an interesting mass range. In this case, we may consider even richer possibilities; for example, ϕ_\pm may be responsible for the 750 GeV excess with an apparent wide width, while η' may be responsible for the 2 TeV diboson excess.

Collider phenomenology of ψ , χ , and φ , is essentially the same as that in the model without the singlet, except that if there are couplings of the form $h^\dagger \Psi_L \bar{\Psi}_N$ and $h \Psi_N \bar{\Psi}_L$, giving a (small) mixing between λ and the Higgs boson, then φ may decay into two Higgs bosons through mixing involving λ . The ξ particle is pair produced by ordinary QCD processes. Its signals depend on the lifetime, and are similar to those of the charge $\pm 1/3$ component of χ . The λ particle is pair produced by electroweak processes. Signals again depend on the lifetime. Suppose it is stable at collider timescales. The electrically charged component of λ is heavier than the neutral component by about 360 MeV. The charged component thus decays into the neutral component and a charged (standard model) pion with a decay length of ~ 10 mm, which might be observed as a disappearing track. Since the decay length is short, however, the current LHC data do not constrain the mass of λ . If λ decays promptly, then it can be observed as a resonance. Since λ has the same standard model gauge quantum numbers as the standard model Higgs doublet, possible decay modes of λ resemble those of heavy Higgs bosons in two Higgs doublet models.

4.2 Hidden quasi-conserved quantities and cosmology

In the model with a singlet hidden quark, approximate symmetries in the G_H sector can be more easily broken than in the model without it. For this purpose, the mass of the singlet hidden quark need not be smaller than Λ , so the following discussion applies even if the mass of the singlet hidden quark is larger than the hidden dynamical scale.

“ L number” can be broken by a renormalizable term

$$\mathcal{L} \sim h^\dagger \Psi_L \bar{\Psi}_N + h \Psi_N \bar{\Psi}_L + \text{h.c.}, \tag{4.14}$$

which mixes λ with the standard model Higgs field.⁸ Unless this mixing is significantly suppressed, λ decays promptly into a pair of standard model particles, with the decay channels similar to heavy Higgs bosons in two Higgs doublet models. Furthermore, if $m_\chi > m_\xi$, χ decays into ξ and a standard model Higgs or gauge boson through the emission of (off-shell) λ . “ D number” is broken by operators

$$\mathcal{L} \sim \left\{ \Psi_D \bar{\Psi}_N \right\} \times \left\{ \begin{matrix} ud \\ ql \end{matrix} \right\} + \text{h.c.} \tag{4.15}$$

These operators induce decay of ξ into a pair of standard model fermions. They also induce decay of χ into a pair of standard model fermions and λ (or a standard model Higgs or gauge boson through the mixing of eq. (4.14)). Since the operator in the first bracket of eq. (4.15) is scalar, it may be significantly enhanced by conformal dynamics of G_H . Note that the dimension of the operators in the second bracket of eq. (4.15) is smaller by one

⁸If G_H is in a conformal phase, the dimensions of operators $\Psi \bar{\Psi}$ are expected to be smaller than 3. In this case, if the coefficients of the interactions in eq. (4.14) are of order one, these interactions participate in the strong dynamics and affect the phase of the theory. Since the effect of this is not clear, here we assume that the coefficients are smaller than order unity at all scales. This requires the coefficients to be much smaller than order one at the scale M_* . This smallness of the coefficients at M_* can be understood by an appropriate chiral symmetry, although this symmetry cannot commute with $SU(5)$ if we want to keep the size of the D number violating interactions in eq. (4.15).

than that of eq. (3.6). Thus, if the G_H sector is in a conformal phase between M_* and Λ , then the required size of anomalous dimensions of $\Psi\bar{\Psi}$ to achieve the same decay rate of χ is smaller by one than in the model without the singlet hidden quark.

Hidden baryon number can also be violated by non-renormalizable interactions involving the singlet hidden quark. For $N = 3$ and $N = 4$, the lowest dimension operators violating hidden baryon number still have the same dimensions as those in eqs. (3.13) and (3.18). For $N = 5$, however, the following operators exist:

$$\mathcal{L} \sim \mathbf{5}^* \Psi_{\mathbf{5}^*} \Psi_{\mathbf{5}^*} \Psi_{\mathbf{5}^*} \Psi_{\mathbf{5}^*} \Psi_N + \mathbf{10} \Psi_{\mathbf{5}^*} \Psi_{\mathbf{5}^*} \Psi_N \Psi_N \Psi_N + \mathbf{10} \bar{\Psi}_{\mathbf{5}^*} \bar{\Psi}_{\mathbf{5}^*} \bar{\Psi}_{\mathbf{5}^*} \bar{\Psi}_N \bar{\Psi}_N + \text{h.c.}, \quad (4.16)$$

in which the dimension of the standard model operators is smaller by one than that in eq. (3.21). Here, to make the expression compact, we have combined the hidden quarks and standard model fermions into $SU(5)$ multiplets: $\mathbf{5}^* = \{d, l\}$, $\mathbf{10} = \{q, u, e\}$, $\Psi_{\mathbf{5}^*} = \{\Psi_D, \Psi_L\}$, and $\bar{\Psi}_{\mathbf{5}^*} = \{\bar{\Psi}_D, \bar{\Psi}_L\}$. These operators induce mixings between hidden baryons and standard model fermions. Because of the lower dimensionality, the required size of anomalous dimensions of the hidden baryonic operators (in the case that G_H is in a conformal phase between M_* and Λ) is smaller by one than in the model without the singlet hidden quark.

With superparticles around the TeV scale, approximate symmetries in the G_H sector are even more easily broken. In particular, D and L numbers are broken by the following superpotential operators

$$W \sim H_u \Psi_L \bar{\Psi}_N + H_d \Psi_N \bar{\Psi}_L + \mathbf{10} \mathbf{10} \Psi_N \bar{\Psi}_{\mathbf{5}^*} + \mathbf{10} \mathbf{5}^* \Psi_{\mathbf{5}^*} \bar{\Psi}_N, \quad (4.17)$$

where H_u and H_d are the up-type and down-type Higgs superfields, respectively. Since these are relatively lower-dimensional operators, particles whose stability would be ensured by D and L numbers, i.e. χ , ξ and λ , decay with cosmologically short lifetimes. For example, with the dimension-five operators suppressed by the scale $M_* \lesssim 10^{17}$ GeV, these particles decay with a lifetime shorter than $O(100 \text{ sec})$ even if G_H is not in a conformal phase between M_* and the TeV scale. With G_H not being in a conformal phase, the coefficients of the first two operators need not be suppressed much: coefficients of $O(0.1)$ or smaller are enough to make the mixing between λ and Higgs fields sufficiently small to preserve their respective phenomenology. Hidden baryon number can also be easily broken. For $N = 3$, there exist dimension-five superpotential operators violating hidden baryon number. With the suppression scale $M_* \lesssim 10^{17}$ GeV, the lifetime of hidden baryons is shorter than $O(100 \text{ sec})$ even if G_H is not in a conformal phase between M_* and the TeV scale. Having conformal dynamics even allows for the decay before the onset of the big bang nucleosynthesis, although the coefficients of the first two operators in eq. (4.17) need to be appropriately suppressed in this case.

5 Discussion

In this paper we have studied a simple model in which the recently reported diphoton excess arises from a composite pseudo Nambu-Goldstone boson, produced by gluon fusion

and decaying into two photons. In the minimal version, the model only has a new hidden gauge group G_H at the TeV scale with a hidden quark in the vectorlike bifundamental representation of G_H and $SU(5) \supset G_{\text{SM}}$. We have found that the model predicts hidden pions $\psi(\mathbf{Adj}, \mathbf{1}, 0)$, $\chi(\square, \square, -5/6)$, and $\varphi(\mathbf{1}, \mathbf{Adj}, 0)$, in addition to the diphoton resonance $\phi(\mathbf{1}, \mathbf{1}, 0)$, and that the masses of ψ and χ are smaller than 1.6 TeV and 1.2 TeV, respectively. The existence of these particles, therefore, can be probed at the LHC in the near future. We have studied physics of would-be stable particles— χ and low-lying hidden baryons—in detail, including constraints from cosmology. We have discussed possible theoretical structures above the TeV scale, including conformal dynamics and supersymmetry, and their phenomenological implications.

In the extended version of the model, there is an additional hidden quark that is singlet under the standard model gauge group and has a mass smaller than the hidden dynamical scale. This yields two hidden pions that can be produced by gluon fusion and decay into standard model dibosons, including a diphoton. We have discussed several scenarios in which these and other resonances can be used to explain various excesses seen in the LHC data. The existence of the singlet hidden quark also helps to write operators inducing decays of would-be stable particles, such as $\chi(\square, \square, -5/6)$, $\xi(\square, \mathbf{1}, -1/3)$, $\lambda(\mathbf{1}, \square, 1/2)$ hidden mesons and low-lying hidden baryons. In particular, if the theory becomes supersymmetric near the TeV scale, the scale suppressing these higher dimensional operators can be as high as the unification scale $M_* \simeq 10^{16}$ GeV while avoiding all the cosmological bounds.

While we have presented it as that explaining the 750 GeV diphoton excess, the model discussed here may also be used to explain other diphoton/diboson excesses that might be seen in future data at the LHC or other future colliders.⁹ In particular, our studies of would-be stable particles and the structure of theories at higher energies can be applied in much wider contexts. We hope that some (if not all) of the analyses in this paper are useful in understanding future data from experiments.

Acknowledgments

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A Hidden sector CP

In this appendix, we consider the effects of possible CP violation in the G_H sector on the standard model. We work in the field basis in which the θ parameter of the G_H gauge theory vanishes: $\theta_H = 0$. (We take this basis by rotating the phase of $\Psi_L \bar{\Psi}_L$, not of $\Psi_D \bar{\Psi}_D$,

⁹We could use the model to explain the 2 TeV diboson excess seen in the 8 TeV ATLAS data [21] by decays of ϕ produced by gluon fusion. This, however, requires $m_\phi \simeq 2$ TeV and $f \simeq 100$ GeV, since the $WW + ZZ$ cross section at 8 TeV is given by $4.7 \text{ fb} (N/5)^2 (100 \text{ GeV}/f)^2$. Therefore, the hidden pion picture is not good, i.e. $m_\phi \gtrsim \Lambda$, in this case.

in order not to affect the QCD θ parameter.) In this case, all the CP violating effects in the G_H sector are encoded in the phases of the hidden quark masses, which we denote with the capital letters, M_D and M_L , to remind ourselves that they are in general complex.

We first analyze the effect on the QCD θ parameter. For this purpose, we consider a nonlinearly realized field

$$U(x) = \exp \left[\frac{2i}{f} \sum_{A=1}^{24} \xi^A(x) t^A \right], \quad (\text{A.1})$$

where t^A are the generators of $SU(5)$, normalized such that $\text{tr}[t^A t^B] = \delta^{AB}/2$, and $\xi^A(x)$ are canonically normalized hidden pion fields. The relevant part of the Lagrangian is then given by

$$\begin{aligned} \mathcal{L} = & -\frac{f^2}{4} \text{tr}[(\mathcal{D}_\mu U)(\mathcal{D}^\mu U)^\dagger] + \frac{c}{2} \text{tr}[MU^\dagger + M^\dagger U] \\ & - \frac{iNg_3^2}{128\pi^2} \text{tr}[(\{t^a, t^b\} + \{t^{a*}, t^{b*}\}) \ln U] \epsilon^{\mu\nu\rho\sigma} G_{\mu\nu}^a G_{\rho\sigma}^b, \end{aligned} \quad (\text{A.2})$$

where the first term is the kinetic term, the second term arises from the hidden quark masses

$$M = \begin{pmatrix} M_D \mathbf{1}_{3 \times 3} & 0 \\ 0 & M_L \mathbf{1}_{2 \times 2} \end{pmatrix}, \quad (\text{A.3})$$

and the third term is determined by chiral anomalies in which we have kept only the gluon field, where t^a ($a = 1, \dots, 8$) are the generators of $SU(3) \subset SU(5)$ corresponding to the standard model color. In our field basis, $c > 0$.

The Lagrangian of eq. (A.2) induces a vacuum expectation value of the ξ^{24} field, which corresponds to the hidden pion ϕ : $\xi^{24} = \langle \xi^{24} \rangle + \phi$. The relevant terms in eq. (A.2) are

$$\mathcal{L} \supset -\frac{1}{2} \partial_\mu \xi^{24} \partial^\mu \xi^{24} - V(\xi^{24}) + \frac{Ng_3^2}{32\sqrt{15}\pi^2 f} \xi^{24} \epsilon^{\mu\nu\rho\sigma} G_{\mu\nu}^a G_{\rho\sigma}^a, \quad (\text{A.4})$$

where

$$V(\xi^{24}) = -3c m_D \cos\left(-\frac{2}{\sqrt{15}f} \xi^{24} + \theta_D\right) - 2c m_L \cos\left(\frac{3}{\sqrt{15}f} \xi^{24} + \theta_L\right). \quad (\text{A.5})$$

Here, we have defined $M_D = m_D e^{i\theta_D}$ and $M_L = m_L e^{i\theta_L}$. The potential of eq. (A.5) gives

$$m_D \sin\left(-\frac{2}{\sqrt{15}f} \langle \xi^{24} \rangle + \theta_D\right) = m_L \cos\left(\frac{3}{\sqrt{15}f} \langle \xi^{24} \rangle + \theta_L\right). \quad (\text{A.6})$$

The solution for $\langle \xi^{24} \rangle$ takes the form

$$\frac{2}{\sqrt{15}f} \langle \xi^{24} \rangle = \theta_D + g(3\theta_D + 2\theta_L; m_D, m_L), \quad (\text{A.7})$$

where $g(x; m_D, m_L)$ is the solution to

$$m_D \sin(-g) = m_L \sin\left(\frac{3}{2}g + \frac{x}{2}\right), \quad (\text{A.8})$$

which can be expanded for small x as $g(x; m_D, m_L) = -(m_L/(2m_D + 3m_L))x + O(x^2)$. If $m_D = m_L = 0$, $\langle \xi^{24} \rangle$ is undetermined at this level, and ξ^{24} would act as the QCD axion.

Through the last term in eq. (A.4), a nonzero value of $\langle \xi^{24} \rangle$ generates a contribution to the QCD θ parameter

$$\Delta\theta = -\frac{2N}{\sqrt{15}f} \langle \xi^{24} \rangle = -N \{ \theta_D + g(3\theta_D + 2\theta_L - \theta_H; m_D, m_L) \}, \quad (\text{A.9})$$

where we have restored θ_H using the fact that only the quantities invariant under the phase rotation of $\Psi_L \bar{\Psi}_L$ can appear here. This expression, therefore, applies in any field basis. In the lack of a low-energy adjustment mechanism such as a QCD axion (and barring accidental cancellation), this contribution must be tiny, $|\Delta\theta| \lesssim 10^{-9}$. Aside from the possibility of fine-tuning between the first and second terms in eq. (A.9), this requires

$$\left| \frac{2m_D\theta_D - 2m_L\theta_L + m_L\theta_H}{2m_D + 3m_L} \right| \lesssim 10^{-9} \frac{1}{N}. \quad (\text{A.10})$$

For $m_{D,L} \neq 0$, this generically forces all the physical phases to be tiny, rendering it unlikely that the hidden η' decays into two hidden pions with a significant rate. We stress that this constraint disappears if there is a QCD axion.

In the presence of a QCD axion, the contribution to θ is harmless. We, however, still expect that the presence of CP violation in the G_H sector generates the following dimension-six operator

$$\mathcal{L} = \frac{1}{\kappa^2} d^{abc} \epsilon^{\nu\rho\sigma\lambda} G_{\mu\nu}^a G_{\rho\sigma}^b G_{\lambda}^{c\mu}. \quad (\text{A.11})$$

This operator induces the neutron electric dipole moment of order $d_n \sim 10^{-26} e \text{ cm} \times (100 \text{ TeV}/\kappa)^{-2}$ [57], so the scale κ must satisfy $\kappa \gtrsim 100 \text{ TeV}$. In the present model, we have

$$\frac{1}{\kappa^2} \sim \frac{g_3^3 N}{16\pi^2} \frac{m}{\Lambda^3} h(3\theta_D + 2\theta_L - \theta_H; m_D, m_L), \quad (\text{A.12})$$

where m collectively denotes $m_{D,L}$, and $h(x; m_D, m_L)$ is a dimensionless function. In the relevant parameter region of $m \approx O(100 \text{ GeV})$, $\Lambda \sim$ a few TeV, $N \sim$ a few, we find $\kappa \sim 100 \text{ TeV}$, so the contribution can be sufficiently small and yet may be seen in future experiments. The G_H sector CP violation also contributes generally to the quark/lepton electric dipole and chromoelectric dipole operators, $\mathcal{L} \sim i\bar{\psi}_L \sigma_{\mu\nu} \psi_R F^{\mu\nu} + \text{h.c.}$ and $i\bar{\psi}_L \sigma_{\mu\nu} t^a \psi_R G^{a\mu\nu} + \text{h.c.}$, but these contributions are suppressed by an extra loop factor, so they are not very constraining. Overall, if there is a QCD axion, we may consider the possibility of significant CP violation in the G_H sector, making the hidden η' decay dominantly into two hidden pions.

We finally mention that CP violation in the hidden sector may generate the operator $\mathcal{L} \sim (g_2^4/(4\pi)^3) m \varphi^\alpha h^\dagger \sigma^\alpha h$ through a loop involving electroweak gauge bosons. This induces a vacuum expectation value of φ : $\langle \varphi^0 \rangle \sim (g_2^4/(4\pi)^3) (m \langle h \rangle^2 / m_\varphi^2) \tilde{h}(3\theta_D + 2\theta_L - \theta_H; m_D, m_L)$, where $\tilde{h}(3\theta_D + 2\theta_L - \theta_H; m_D, m_L)$ is a dimensionless function. The induced vacuum expectation value, however, is too small, $\langle \varphi^0 \rangle \lesssim 10^{-3} \text{ GeV}$, to affect the phenomenology of hidden pions or the electroweak precision measurements.

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