Challenges for models with composite states

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Composite states of electrically charged and QCD-colored hyperquarks (HQs) in a confining SU(N_{HC}) hypercolor gauge sector are a plausible extension of the standard model at the TeV scale, and have been widely considered as an explanation for the tentative LHC diphoton excess. Additional new physics is required to avoid a stable charged hyperbaryon in such theories. We classify renormalizable models allowing the decay of this unwanted relic directly into standard model states, showing that they are significantly restricted if the new scalar states needed for UV completion are at the TeV scale. Alternatively, if hyperbaryon number is conserved, the charged relic can decay into a neutral hyperbaryon. Such theories are strongly constrained by direct detection, if the neutral constituent hyperquark carries color or weak isospin, and by LHC searches for leptoquarks if it is a color singlet. We show that the neutral hyperbaryon can have the observed relic abundance if the confinement scale and the hyperquark mass are above TeV scale, even in the absence of any hyperbaryon asymmetry.

I. INTRODUCTION

Hints of a small excess of events in the diphoton channel have been reported by ATLAS [1, 2] and CMS [3, 4] experiments during the 13 TeV run. Although they are probably a statistical fluctuation, it is intriguing that both experiments see the excess at the same invariant mass of the photon pair, at approximately $m_{\gamma\gamma} = 750$ GeV, and that there are further hints of an excess at higher invariant masses [2]. This has prompted numerous theoretical interpretations in terms of a spin-0 resonance decaying into photons. A plausible class of models considers the resonance to be composed of heavy quarks (HQs), bound by an SU(N_{HC}) hypercolor confining gauge theory. If the HQs carry electric charge then the pion- or quarkonium-like bound state assumed to be the 750 GeV resonance can decay into photons to explain the observed signal [5, 17].

This kind of extension of the standard model seems natural since it simply enlarges the gauge symmetry group by an additional SU(N_{HC}) factor. One might expect that, similarly to the standard model, the HQs carry a new, possibly conserved charge, hyperbaryon (HB) number. In this work we explore the consequences of HB number being conserved, leading to a dark matter candidate but also potentially severe conflicts with observation.

1 For a discussion of more exotic hyperbaryons, consisting of bound states of hyperquarks in several different representations of the standard model gauge symmetries, see ref. [18].

We will assume that $\Psi$ is vectorlike. If $m_{\Psi} \ll \Lambda_{HC}$ then the bound state $\tilde{\Psi}$ is pion-like, with a mass scaling as $m_{\tilde{\Psi}} \sim \sqrt{m_{\Psi} \Lambda_{HC}}$ due to an assumed approximate chiral symmetry, softly broken by $m_{\Psi}$. If $m_{\Psi} \gg \Lambda_{HC}$, the composite state would be more similar to charmonium. However the fact that the putative 750 GeV resonance is relatively narrow and distinct indicates that $m_{\Psi}$ cannot be much greater than $\Lambda_{HC}$; otherwise one would expect to produce a series of closely-spaced resonances with fractional mass splitting $\Delta m/m \sim (\Lambda_{HC}/m_{\Psi})^{3/4}$, based upon a semiclassical model of bound states in a linear confining potential $V \sim \Lambda_{HC}^2 r$.

In these models, it is assumed that the HQs are also colored and can thus be produced by gluon-gluon fusion ($ggF$). Alternatively, it is possible to have sufficient production through photon-photon or vector boson fusion [19, 23] if the HQ carries a large hypercharge $\sim (3-4)$, indicating a Landau pole at a relatively low scale, barring additional states. However the modest growth of the photon parton distribution function (PDF) with energy puts this scenario in tension with the lack of any ob-
served signal in the 8 TeV LHC run. For these reasons models with $ggF$ production are favored, and our focus will therefore be on HQs that carry QCD color. We will comment on colorless models in section VI.

One possibility is that there is a conserved HQ number that leads to a stable hyperbaryon (HB) consisting of $N_{HC}$ HQs. There are very stringent constraints on electrically charged relics, so that a realistic model should provide some way for these unwanted relics to decay. In some cases it is possible to write down a high-dimensional effective operator that would allow the charged HB to decay directly into standard model (SM) particles. However it is theoretically more satisfying to demonstrate the renormalizable interactions that would allow for the decays of the HB, either into purely SM particles, or into an electrically neutral HB that might be a viable dark matter candidate. One point of the present work is that there are relatively few categories of renormalizable models that lead to nonconserved HB number, and they are strongly constrained by collider searches if the new scalars that lead to a stable hyperbaryon (HB) consisting of $N_{HC}$ HQs. There are very stringent constraints on electrically neutral HQs. We survey these possibilities in section II.

If HB number is conserved, the lightest HB must be electrically neutral, and is a dark matter candidate. A priori, it could carry SU(3)$_c$ color, in which case it binds to ordinary quarks or antiquarks to make a color-neutral composite state, whose residual strong interactions give it a large cross section for scattering on nucleons. This is likely to be excluded at the same level as charged relics by searches for anomalous heavy isotopes, favoring models in which the lightest stable HQ is a color singlet. If it is an SU(2)$_L$ doublet, the constraints from direct detection are less severe, but still quite significant. The safest case with constituents that are purely neutral under SM interactions turns out to have a relatively light leptoquark bound state, assuming that ordinary baryon number is still an accidental symmetry of the full theory; hence even this case comes under pressure from current LHC constraints. These scenarios are discussed in section III.

The relic density of conserved HBs could be due to an asymmetry, analogous to the baryon asymmetry. Here we suppose that the mechanism for generating an HB asymmetry is weak or lacking, and focus on the symmetric component, which should be understood in any case before invoking an asymmetry. In section III we compute the abundance for purely singlet HBs and for those that carry weak isospin, showing its dependence on the confinement scale $\Lambda_{HC}$ and the mass of the neutral HQ. Even the purely singlet HB has electromagnetic interactions with protons through loops containing the charged HQ. This leads to weak constraints from direct detection that we derive in section IV. Models in which the charged HQ does not carry QCD color are much less restricted by the considerations of the previous sections. We briefly comment on them in section VI and give conclusions in section VII.

II. HYPERBARYON NUMBER VIOLATION

First we deal with the possibility that hyperbaryon number is not a symmetry of the theory, and HBs can decay directly into SM particles. One might imagine further possibilities by allowing new dark matter particles in the final states, but we do not pursue this here.

Let us provisionally assume that the charged HQ is a fermion $\Psi^A_i$ with HC index $A$, color index $a$, and weak hypercharge $Y$. In addition, there may be bosonic fields carrying fundamental HC indices. We can associate the global HB quantum number 1 to each fundamental HC index, and −1 to each antifundamental index. Thus a HC gauge boson, having one of each, has vanishing HB number. If we were only allowed to contract HC indices in fundamental/antifundamental pairs, then it would be impossible to violate HB number while respecting the gauge symmetry. However in SU(N) we also have the invariant tensor $\epsilon_{A_1 \ldots A_N}$. This simple argument demonstrates that HB violation must involve the $\epsilon$ tensor.

We start by discussing a rather general class of models that lead to decays of the hyperbaryons into standard model quarks. Other types of models have a structure depending upon the value of $N_{HC}$, so we consider the possibilities $N_{HC} = 2, 3, 4$ in turn in the following.

II.1. Neutral scalar hyperquarks

There is a general class of models which have the same structure and are renormalizable for $N_{HC} = 2, 3, 4$, requiring the presence of $N_{HC}$ flavors of fundamental SU($N_{HC}$) scalars $\Phi_i, A$, and that the hypercharge of $\Psi$ matches that of the SM $u_R$ or $d_R$ quarks. Illustrating the former case where $Y_{\Psi} = 2/3$, it has the form

$$\sum_{i=1,j} \lambda_{ij} \bar{\Psi}^A_i \Phi_i, A u^0_{R,j} + \mu \epsilon_{A_1 \ldots A_N} \Phi_{1, A_1} \cdots \Phi_{N, A_N}$$

(1)

where $\mu$ has dimensions of $(mass)^{4-N}$ and $N = N_{HC}$. If the $\Phi$'s are heavy they can be integrated out giving an operator schematically of the form

$$\frac{\mu \lambda^N}{m_{\Psi}^2} (\bar{\Psi} u_R)^N$$

(2)

that allows the charged HB to decay into $N$ up-type quarks.

On the other hand if the $\Phi$'s are lighter than $m_{\Psi}$, then the lightest HB is a scalar bound state $\Phi^N$, which can decay via the operator $\mu \Phi^N$. Concretely, this would allow the HB to decay into $N - 1$ hypermesons,

$$\Phi^N \rightarrow (N - 1) \Phi^* \Phi$$

(3)
LHC searches for heavy quarks. There will be Drell-Yan mixing comes from the Yukawa coupling $\lambda$ between the SM quarks and the composite fermions $\Psi$. If final state $\Phi$ annihilate with $\Phi$ before hadronizing). If $N \geq 4$ there is generically enough phase space for the decay, even if the masses are dominated by the constituents, since $N m_\Phi \geq 4 m_\Phi$. For $N = 3$ one would require the mesons to be pseudo-Goldstone bosons in order to overcome this restriction. (For $N_{HC} = 2$ there is no clear distinction between a meson and a baryon since the fundamental representation of SU(2) is pseudoreal.)

II.1.1. LHC constraints

If $\Phi$ is heavy, then in addition to the HB-violating operator (2), there is a dimension-6 HB-conserving interaction of the form

$$\frac{\lambda^2}{m_\Phi^2} |\bar{\Psi} \gamma_5 U_R|^2$$

Because of its chiral structure it would allow vector hypermesons to decay into $u_R$ quarks plus a gluon, leading to three jets [1]. However there is no obstruction to making $\lambda^2 / m_\Phi^2$ sufficiently small so that the branching ratio for these decays is unimportant. A lower bound (depending upon $N_{HC}$ and $\mu$) can be placed on $\lambda / m_\Phi^2$ by demanding that the charged HB decays become before big bang nucleosynthesis (BBN), leading to $\lambda / m_\Phi^2 \gtrsim 10^{-4} / \text{TeV}^2$ for the most restrictive case of $N_{HC} = 4$, consistent with a small coefficient $\sim 10^{-8} / \text{TeV}^2$ for the effective operator (4), if $m_\Phi \sim \text{TeV}$.

The case of light $\Phi$ is more interesting, since it leads to mass mixing between the SM quarks and the composite fermions $U_i \equiv \bar{\Psi}_i \Phi_i$, that have the same number quantities as $u_R$ (or $d_R$ in the alternate $\Psi = -1/3$ models). The mixing comes from the Yukawa coupling $\lambda_{ij} \bar{\Psi}_i \Phi_j u_{R,j}$ leading to an off-diagonal mass term of order $\lambda_{ij} f_\Phi U_i u_{j}$, where $f_\Phi$ is the hypermeson decay constant.

The heavy composite $U_i$ particles are constrained by LHC searches for heavy quarks. There will be Drell-Yan pair production of $\bar{U}_i U_i$, followed by decays $U_i \to u_i h$ where $h$ is the Higgs boson, or $U_i \to W d_i$. The first decay comes from mixing of $U_R$ with $u_R$, while the second is due to $U_L-u_L$, mixing. We note that $U$ is a Dirac fermion since we have implicitly assumed that $\Psi$ is vectorlike in order to have a bare mass. The mass matrix takes the form

$$(\bar{u}_R U_R) \begin{pmatrix} m_u & \lambda f_\Phi \mu \\ 0 & M_U \end{pmatrix} \begin{pmatrix} u_L \\ U_L \end{pmatrix}$$

where $m_u$ is the SM quark mass matrix and $M_U \sim \delta_{ij} \sqrt{\Lambda_{HC}} m_\Psi \lesssim 750 \text{ GeV}$ is the mass matrix of the composite states in the absence of $U-u$ mixing. After diagonalization, the left- and right-handed states have mixing angles

$$\theta_L \sim \frac{\lambda f_\Phi m_u}{M_U}, \quad \theta_R \sim \frac{\lambda f_\Phi}{M_U}$$

The effective couplings for $U \to h u$ and $U \to W d$ are thus of order

$$\frac{m_u \theta_R}{v} = \frac{\lambda \Lambda m_u}{v M_U}, \quad g_{2 \theta L} = \frac{g_2 \lambda \Lambda m_u}{M_U}$$

implying that $U \to h u$ is the dominant decay channel, as long as $m_U \gtrsim 300 \text{ GeV}$.

This scenario has been considered by ATLAS [24] and CMS [25] (see also [26]) for the case of top partners decaying as $T \to h t$ and $W b$. These searches constrain $m_T > 700 - 900 \text{ GeV}$ depending upon the respective branching ratios, with the strongest limit when $T \to h t$ dominates, as we expect here. This contradicts the assumption that $m_\Psi < m_\Phi$ which would imply that $m_T < 750 \text{ GeV}$, hence ruling out a dominant coupling to top quarks.

II.1.2. Flavor constraints

The interaction (1) induces flavor changing neutral current decays $c \to u \gamma$ as shown in fig. 1. The analogous diagram gives $b \to s \gamma$ in the related model with $\tilde{Y}_\Psi = -1 / 3$. Defining $t = m_{\Psi}^2 / m_{\tilde{\Psi}}^2$, and writing the transition amplitude as $(m_b / \Lambda_b^2) \tilde{s}_R \sigma_{\mu\nu} q_\nu b_L$, we find [27]

$$\frac{1}{\Lambda_b^2} = \frac{N_{HC} (e / 3) \lambda_6 \lambda_8}{32 \pi^2 m_\Phi^2} f(t) < \left( \frac{1}{55 \text{ TeV}} \right)^2$$

where $f(t) = t / (t - 1)^4 [(t - 1)(t^2 - 5 t - 2) / 6 + t \ln t] \sim 0.1$ for a typical value $t \sim 1.5$ and the experimental upper limit is inferred from ref. [28] (specifically, $\Lambda_b^{-2} = 4 \sqrt{2} G_F V_{tb} V_{tb}^* C_7 e / 16 \pi^2$ with $C_7 < 0.065$). Taking for example $m_\Psi = 400 \text{ GeV}$ and $N_{HC} = 3$ we obtain a weak constraint, $\sqrt{|\lambda_6 \lambda_8|} \lesssim 0.75$.

II.2. $N_{HC} = 2$

We turn next to models that are specific to the value of $N_{HC}$. For $N_{HC} = 2$ we can construct the HB-violating
dimension-6 operators

\[ \epsilon_{abc} \epsilon_{AB}(\overline{\Psi}_A \Psi_B) \left( \overline{u}_{R}^{c}l_{R}, \overline{E}_{L}^{c}, d_{R}^{c}l_{R} \right) \]

where SU(2) \_L indices are implicitly contracted with \( \epsilon_{ab} \) in the second operator. The first two require the electric charge of \( \Psi \) to be \( q_{\Psi} = 1/6 \) while the last one needs \( q_{\Psi} = 2/3 \). They can be UV-completed by introducing a color-triplet scalar with couplings

\[ \epsilon_{abc} \epsilon_{AB}(\overline{\Psi}_A \Psi_B) \Phi_c + \Phi^a (\overline{u}_{R}^{c}l_{R}, \overline{E}_{L}^{c}, d_{R}^{c}l_{R}) \]

(we omit writing the dimensionless coupling constants). The \( \overline{\Psi} \Phi \) interaction makes it clear that there is a bound state \( \Phi = \overline{\Psi} \Psi \) with the same quantum numbers as \( \Phi \). The \( \overline{\Psi} \Phi \) operator becomes an off-diagonal mass term \( \sim f_{\Psi}^2 \Phi \Phi \) (where \( f_{\Psi} \) is the hypermeson decay constant) that causes mixing between the elementary and composite scalars. Therefore the experimental constraints on leptoquarks also apply to \( \Phi \), assuming its production cross section is the same as that of \( \Phi \).

Generally, the production of \( \Phi \) will be of the same order as that for \( \Psi \), depending upon the probability for \( \overline{\Psi} \Psi \) to hadronize into \( \Phi \Phi \) versus other hadron-like pairs. In the present case, there are only two ways to hadronize, either into mesons \( \overline{\Psi} \Psi \) or baryons \( \overline{\Psi} \Psi \), so we expect the production cross section to be about half of that for an elementary color triplet pair. Taking this into account, we can infer the CMS limits on the \( \Phi \) mass to be \( m_{\Phi} \gtrsim 680 \text{ GeV} \) if \( \Phi \to Tb \) predominantly and \( m_{\Phi} \gtrsim 650 \text{ GeV} \) if \( \Phi \to \mu q \). These are marginally compatible with the expected value \( m_{\Phi} \sim 750 \text{ GeV} \). The corresponding ATLAS limits are similar.

II.3. \( N_{hc} = 3 \)

If \( N_{hc} = 3 \) and \( \Psi \) carries hypercharge \( Y = 1 \), it can couple to right-handed leptons \( e_{R,i} \) of generation \( i \), and a neutral colored scalar hyperquark \( \Phi \),

\[ \lambda_i \overline{\psi}_A^i \Phi_A e_{R,i} + \mu \epsilon^{ABC} \epsilon_{abc} \Phi_A^a \Phi_B^b \Phi_C^c \]

(11)

Supposing that the scalar is heavy, one can integrate it out to obtain a dimension-9 operator schematically of the form

\[ \frac{\mu \lambda^3}{m_{\Phi}^3}(\overline{\Psi} \Psi e_{R}) \]

(12)

that allows the charge-3 relic HB to decay into three leptons. This effective operator was pointed out in ref. [9], where \( \Psi \) carried an extra flavor index \( f = 1, 2 \), necessitating the existence of all possible combinations of \( \Psi_1 \) and \( \Psi_2 \) in operators like (12) to deplete all the flavors of baryons. We note that the UV-completion solves another problem of their model, namely the overabundance of hyperons of the form \( \overline{\pi}_{12} = \overline{\Psi}_1 \Psi_2 \) that were stable in the theory with only the effective operator (12), but become unstable to \( \overline{\pi}_{12} \to e_{R} \tilde{e}_{R} \) by \( \Phi \) exchange using interactions of the type (11).

II.3.1. Constraints on light \( \Phi \)

If \( \Phi \) is relatively light, then analogously to the discussion in section II.3, there will be a vector-like composite state \( E = \overline{\Psi} \Psi \) that has the same quantum numbers as \( e_{R,i} \), and we get mass mixing between \( E \) and a linear combination of the SM leptons \( e_{R,i} \). ATLAS has searched for the decays of a vector-like lepton into \( Z \) and a SM lepton \( Z \). The constraints are not very restrictive, ruling out the mass ranges 129-176 GeV (114-168 GeV) if the mixing is primarily to electrons (muons), except for gaps 144-163 GeV (153-160 GeV) where a heavy lepton is still allowed.

It is also possible to derive constraints from the rare flavor-violating decays \( Z \to \ell_i^+ \ell_j^- \). These arise because the unitary transformations \( U_L \) and \( U_R \) that diagonalize the \( 4 \times 4 \) Dirac mass matrices do not act unitarily in the \( 3 \times 3 \) subspace involving only the SM leptons, which couple to \( Z \) whereas the heavy \( E \) state does not. These flavor-changing couplings are proportional to

\[ L_Z = -\frac{g}{2c_w} e_i \left[ (U_{L}^{1} P_{3} U_{R})_{ij} (-1 + 2 s_{W}^2) P_L \right] \]

\[ + (U_{R}^{1} P_{3} U_{L})_{ij} (2 s_{W}^2) P_R e_j \]

(13)

where \( c_w, s_w \) are the weak mixing factors and \( P_3 \) projects onto the \( 3 \times 3 \) subspace of the SM leptons. Taking \( P_3 = 1 - P_4 \) where \( P_4 \) projects onto the heavy state, one can express the nonstandard contributions to (13) as

\[ \delta L_Z = \frac{g f_{\lambda}^2}{2c_w m_E} \overline{\ell}_{i} \ell_{j} \left[ (m_{\ell} \lambda)_{i}(m_{\ell} \lambda)_{j} (-1 + 2 s_{W}^2) P_L \right] \]

\[ + \lambda_{i} \lambda_{j} (2 s_{W}^2) P_R e_j \]

(14)

where \( m_{\ell} \) is the (diagonal) light lepton mass matrix and \( m_E \) is the composite state mass. Because of the \( m_{\ell} \)-suppression of the \( U_{L}^{1} \), mixing, the right-handed couplings dominate. Assuming that \( m_{E} \sim f_{\lambda} \), the experimental upper limits on decays into \( e_{\mu}, e_{\tau}, \mu_{\tau} \) lead to \( \sqrt{\lambda_{\mu} \lambda_{\tau}} \lesssim 0.08, \sqrt{\lambda_{\mu} \lambda_{\tau}} \lesssim 0.14 \).

Somewhat stronger bounds arise from the diagonal contributions, which can induce flavor nonuniversality in flavor-conserving decays \( Z \to \ell_i^+ \ell_j^- \) through interference with the SM amplitudes. Ignoring the small contribution from the left-handed couplings \( g_{L} \), and assuming that \( \lambda_{i} \gg \lambda_{i} \), the fractional deviation \( \Delta R_{\tau/e} = \text{BR}(Z \to \tau \tau)/\text{BR}(Z \to e) - 1 \) is given by

\[ \Delta R_{\tau/e} = 2 \frac{g_{R} \delta g_{R}}{g_{R}^2 + g_{L}^2} = -\frac{8 s_{W}^2}{1 - 4 s_{W}^2 + 8 s_{W}^4} \frac{\lambda_{i}^2 f_{\lambda}^2}{m_{E}^2} \]

(15)
The experimental limit (at 1σ) is $\Delta R_{T/\gamma} > -0.0013$, implying $|\lambda_\gamma| \lesssim 0.04$, again assuming that $m_\psi \sim f_\pi$.

Products $\lambda_i \lambda_j$ with $i \neq j$ are constrained by radiative flavor violating decays at one loop, illustrated by $\mu \to e\gamma$ in fig. [1]. Defining $t = m_\phi^2 / m_\mu^2$ and writing the transition amplitude as $(m_\mu / \sqrt{2}) \sigma^\mu q_\mu \epsilon_\mu \epsilon_L$, we find [27]

$$\frac{1}{N_\mu} = \frac{9e\lambda_\mu \lambda_\mu}{32\pi^2 m_\phi^2} f(t) < \left( \frac{1}{64 \text{ TeV}} \right)^2$$

where $f(t)$ is as in eq. [8] and the experimental limit is inferred from refs. [34, 35]. Taking for example $m_\phi \approx 400$ GeV and $f \sim 0.1$ leads to the limit $|\lambda_\mu \lambda_\mu| < 0.2$, less restrictive than that from $Z \to \mu e$.

II.4. $N_{hc} = 4$

For $N_{hc} = 4$, it is also possible to violate HB with renormalizable interactions if there exists a colored scalar $\Phi_{AB,a}$ in the antisymmetric tensor representation of $SU(N_{hc})$, as well as a color-triplet fundamental $\Phi_{a}$. One can then construct the interactions

$$\epsilon_{abcd} \epsilon_{A,a} \Phi_{B,b,c} \Phi_{CD,c} + \epsilon_{abcd} \Phi_{CD,a} \Phi_{b} \Phi_{D}$$

$$\epsilon_{abcd} \Phi_{c} \Phi_{D} \Phi_{ab}$$

where $q_R$ can be either $u_R$ or $d_R$. Integrating out the scalars gives the dimension-9 operator

$$\frac{\mu}{M_\Phi^4} \epsilon_{abcd} \epsilon_{efg} (\Phi_{A,a} \Phi_{B,b,c}) (\Phi_{C,c} q_R,d) (\Phi_{D,e} q_R,f)$$

It allows four $\Psi$'s to decay into two quarks, and would mediate decay of the charged hyperbaryon $\Psi^4$, provided that $\Psi$ has charge $1/3$ or $-1/6$. The collider phenomenology stemming from the $\Phi q_R$ operator is similar to that of the model given by eq. [1].

II.5. Summary

We have presented several renormalizable frameworks allowing for deletion of the unwanted charged relic hyperbaryon. Most of them require a charge-neutral scalar hyperquark $\Phi$, that might also be colored if $N_{hc} = 3$. If $\Phi$ is sufficiently heavy to avoid being produced at the LHC, these models are practically unconstrained. On the other hand, if $m_\Phi \lesssim \text{TeV}$, $\Phi$ can form a meson-like bound state with $\Psi$ that has the quantum numbers as a SM quark or lepton. We showed that the first case is rather strongly constrained by ATLAS and CMS searches for heavy quarks (top partners).

If $N_{hc} = 2$, another possibility is to introduce a scalar leptoquark $\Phi$ that is neutral under $SU(N_{hc})$. In this model, there is a bound state of $\Psi\Psi$ with the same quantum numbers that mixes with $\Phi$ and thus introduces a relatively light leptoquark state. This is strongly constrained by ATLAS and CMS, leaving little room for a model in which the $\bar{\Psi}\Psi$ hypermeson could be as light as 750 GeV.

These models typically also predict some level of quark or lepton flavor violation, but we find that the resulting constraints are typically weak. The new Yukawa couplings appearing in $\lambda_i \Phi f_i$, where $f_i$ is a SM fermion, need only be of order 0.1 in most cases. Flavor universality of $Z \to \ell_i \ell_i$ decays gives the strongest such limit, $\lambda_\tau < 0.04$.

III. NEUTRAL HYPERBARYONS

A second possibility is that the charged HQ can decay into a lighter neutral HQ, which we denote by $S^A$, that is fundamental under $SU(N_{hc})$ and electrically neutral. In principle it could also carry QCD color or weak isospin, but as we will show, these options are generally disfavored by constraints from direct detection of the resulting hyperbaryon $S^{N_{hc}}$. We introduce a scalar $\Phi$ that mediates the decay of $\Psi$ to $S$ plus standard model particles through an interaction of the form

$$\lambda S^A \Phi \Psi_a$$

followed by decay of the mediator $\Phi$ into standard model particles. (Indices corresponding to any additional quantum numbers are suppressed here). Alternatively, $\Psi$ and $S$ could be in an SU(2)$_L$ doublet, so that $\Psi \to SW$ by the SM weak interaction. In the following, we consider the different possible cases for additional quantum numbers carried by $S$.

III.1. Colored stable hyperquark

We first consider the case in which the neutral HQ $S$ is colored. The baryonic state that is a singlet under $SU(N_{hc})$ is not color-neutral, if $S$ is fermionic. The $SU(N_{hc})$ singlet operator $\epsilon_{A_1 \ldots A_N} S_{A_1 \ldots A_N}$ is symmetric under interchange of $SU(3)$ indices, and can only be antisymmetric under spin if $N_{hc} = 2$. To make a color singlet, it must bind with ordinary quarks,

$$B = \epsilon_{A_1 \ldots A_N} S_{A_1 \ldots A_N} q_{a_1} \ldots q_{a_N}$$

(20)

whose flavors and spins (or spatial configurations) are chosen so as to make the $q \cdots q$ part of the wave function totally antisymmetric, while maintaining charge neutrality. For example if $N_{hc} = 3$, one can form the antisymmetric $s$-wave state of two down quarks and one up quark, whose flavor/spin wave function is

$$udd (\uparrow\uparrow\downarrow - \uparrow\downarrow\uparrow) + ddu (\uparrow\uparrow\downarrow - \uparrow\downarrow\uparrow) + dud (\uparrow\downarrow\uparrow - \uparrow\uparrow\downarrow)$$

(21)

This bound state of $SSSudd$ could be expected to behave similarly to a heavy neutron in its scattering on ordinary baryonic matter.

The scattering properties of dark matter comprised of exotic baryonic-like bound states have been discussed in
An exception is when \( N_{HC} = 3 \) with colored scalar HQs, denoted by \( \Phi \). In that case the bosonic HB state

\[
B = \epsilon_{ABC} \epsilon_{abc} \Phi^{A,a} \Phi^{B,b} \Phi^{C,c}
\]

is neutral under all gauge symmetries and has the correct statistics. However this model does not have HB conservation as an accidental symmetry, since the super-renormalizable operator (24) can simply appear in the Lagrangian, like in the model of eq. (11). We therefore consider it to be unnatural as an example of HB conservation.

### III.2. Weakly interacting stable hyperquark

In models where the SM gauge indices of the hyperquarks are embedded in the fundamental of SU(5) for gauge unification [5–7, 11, 16, 17], the colored HQ can be expected to decay into a doublet HQ by exchange of a heavy GUT gauge boson. The charged component of the doublet can then decay into the neutral \( S \) through weak interactions. In this case, the hyperbaryon \( S_{N_{HC}} \) will have weak interactions with ordinary matter through \( Z \) exchange, similar to a hypothetical heavy bound state containing \( N \) left-handed neutrinos. In terms of previously studied models, we expect the cross section for scattering on nucleons to be similar to that of Dirac Higgsino dark matter, neglecting the Higgs exchange contributions to the scattering that occur in that model and focusing only on \( Z \) exchange. This has been studied in ref. [11], which finds that the spin-independent cross section for scattering on neutrons is

\[
\sigma \cong 1.0 \times 10^{-37} \text{ cm}^2
\]

This is \( \sim (7–9) \) orders of magnitude above the current LUX limit [42, 43], requiring that the relic density of such HBs be correspondingly depleted. The cross section for HB scattering is expected to be \( N_{HC}^2 \) times larger due to the number of constituents. We show the limit on the fractional abundance as a function of the HB mass in fig. 2, including this dependence on \( N_{HC}^2 \).

### III.3. Singlet stable hyperquark

An interesting possibility is that the scalar mediator carries away the charge and color of the \( \Psi \) hyperquark, appearing in place of \( Q \). For example with \( N_{HC} = 2 \), one can form the state

\[
B_{\Psi} = \epsilon_{AB} \bar{\Psi}^{A,a} \Phi^{B,b} u_{\alpha} u_{\beta} (\uparrow \downarrow - \downarrow \uparrow)
\]

provided the charges of \( \Psi \) and \( u \) are equal. However regardless of these details, we expect that any bound state containing ordinary quarks will bind to protons, and the previous result will hold.

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4 To generalize the previous example to other values of \( N_{HC} \), one must admit nonzero values of the HQ electric charge since charge-neutral combinations of \( N_{HC} \) ordinary quarks do not generally exist. In this case it may be possible to dispense with \( Q \) altogether and find a neutral bound state of the form (25) with \( \Psi \).
leaving $S^4$ charged only under $SU(N_{vc})$. Then the possible couplings of $\Phi$ allowing decays to SM particles, while maintaining ordinary baryon number as an accidental symmetry, are limited to the first two cases shown in table I where the hypercharge $y_{\Phi}$ can take values 7/6 or 1/6. The last two, with $y_{\Phi} = -1/3$ or $-4/3$, are disfavored because they allow for baryon violation by marginal operators, leading to rapid proton decay. In the favored models, baryon number can be consistently assigned to all fields, such that $B$ coincides with HB number.

In the preferred models with $y_{\Phi} = 7/6, 1/6$, the scalar $\Phi$ (and therefore $\Psi$) is an $SU(2)_L$ doublet. These models are viable for producing the 750 GeV diphoton signal in the pion-like regime, for which the ratio of branching ratios $R_{BR} = \text{BR}(\tilde{\tau} \rightarrow \gamma\gamma)/\text{BR}(\tilde{\tau} \rightarrow gg)$ shown in the last column of table I is relevant; this ratio is computed following ref. [3]. If $R$ is too small, the observed diphoton rate would require the width for decays into gluons to be so large (in order to compensate for the small BR into photons) that they would exceed the ATLAS bound on dijets [41], $\sigma(\tilde{\tau} \rightarrow gg) < 2.5\, \text{pb}$ at $\sqrt{s} = 8\, \text{TeV}$. We find the lower limit $R > 1.6 \times 10^{-4}$ to satisfy this constraint; details are given in appendix A. All the models in table I are consistent with this constraint.

However these models suffer from another constraint, namely searches for leptoquarks at the LHC. Even though $\Phi$ may be very heavy, it mixes with bound states of $\bar{\Psi} = \bar{\Psi} S$ that have the same quantum numbers as $\Phi$. These hypermesons will be pair-produced at LHC and their masses must be less than 750 GeV since $m_S < m_{\Psi}$; $m_{\tilde{\Phi}}$ can only be decreased by mixing with $\Phi$. They can decay only into quarks and leptons since they have the quantum numbers of leptoquarks. ATLAS and CMS find lower bounds on scalar leptoquarks decaying into jets and electrons or muons such that $m_{\tilde{\Phi}} \gtrsim 1\, \text{TeV}$ for branching ratio $\beta = 100\%$ into one of those channels, and $m_{\tilde{\Phi}} \gtrsim 800\, \text{GeV}$ for $\beta = 30\%$ [29–32]. Even if $\Phi$ decays mostly into $\tau$ and 3rd generation quarks, the limit ranges from $m_{\tilde{\Phi}} > 500 – 740\, \text{GeV}$ for $\beta = 50 – 100\%$, from the Run I data. Thus there is very little parameter space in which to hide an expected leptoquark with $m_{\tilde{\Phi}} < 750\, \text{GeV}$, making it difficult to accommodate these models.

A conceivable way out might be to choose small dimensionless couplings for (19) and the interactions of table I such that the composite leptoquark is metastable and decays outside of the detector (but with a lifetime still below 1 s to avoid problems with BBN). CMS has searched for such long-lived charged particles [6] Ref. [45] from run

| $y_{\Phi}$ | $T_{3,\Phi}$ | $q_{\Phi}$ | $L_{B}$ | $L_\beta$ | $\text{BR}(\tilde{\tau} \rightarrow \gamma\gamma)/\text{BR}(\tilde{\tau} \rightarrow gg)$ |
|----------|--------------|------------|--------|---------|------------------------------------------|
| $+7/6$   | $\pm 1/2$    | $2/3, 5/3$ | $\{\Phi_u Q^L_u l_R, \Phi_d d_R L^c_R\}$ | none | 0.12                                      |
| $+1/6$   | $\pm 1/3, 2/3$ | $\Phi_u d_R L^c_R$ | none | $2.7 \times 10^{-3}$ |
| $-1/3$   | $0$           | $-1/3$     | $\Phi_u R^{-1} R^{-1}$ | $\{\Phi^* Q^L_u Q^L_{dR}, \Phi^* u_R d_R, \Phi^* u_R u_R\}$ | $4.3 \times 10^{-4}$ |
| $-4/3$   | $0$           | $-4/3$     | $\Phi_d R^{-1} R^{-1}$ | none | 0.11                                      |

TABLE I. Possible hypercharges, weak isospin and electric charges of the colored scalar mediator $\Phi$, the baryon-conserving operators leading to $\Phi$ decay into SM particles, and allowed baryon-violating operators. The last column is the ratio of branching ratios of a pion-like 750 GeV state into photons versus gluons, for constituents $\Psi$ having the same SM quantum numbers as $\Phi$.  

5 CMS reports a slight excess of $eejj$ events corresponding to $m_{\tilde{\Phi}} = 650$, $\beta = 0.015$.  
6 similar searches by ATLAS are difficult to interpret in the context of the present model.
FIG. 4. Inverse Bohr radius of the hyperbaryon bound state versus \( \Lambda_{HC} \), for several values of the HQ mass \( m_S = 3000, 300, 100, 10, 1 \) GeV and \( N = 2, 3, 4 \) (smaller \( N \) gives lower curve for a given mass.) Vertical bars show where the perturbative treatment breaks down and extrapolation of low-\( \Lambda \) behavior is used, for a given \( m_S \).

I directly constrains charged hypermesons, which however are assumed to be produced by Drell-Yan rather than \( gg \) fusion. Ref. \[46\] does a similar analysis in run II, considering DY-produced particles of charge 1 and 2, obtaining limits on the production cross section that we reproduce in Fig. 3. Our prediction for the production of \( \Phi \Phi^* \) pairs (assuming they originate from \( gg \rightarrow \Psi \Psi \) and \( qq \rightarrow \Psi \Psi \) that hadronize mainly into \( \Phi \Phi^* \)) is also shown there.\(^7\) We expect the leptoquark states of charge 2/3 and 5/3 to be constrained at a similar level to the charge 1 and 2 particles considered in \[46\], leading to a limit of \( m_S \gtrsim 1.35 \) TeV, again in contradiction to the premises of the present model.

IV. RELIC DENSITY OF NEUTRAL HYPERBARYONS

If hyperbaryon number is conserved and results in a HB that is neutral under SM gauge interactions, it could be a viable dark matter candidate. If it carries weak isospin, then the considerations of section III.2 show that it can only be a very subdominant component of the total dark matter. In either case, it is interesting to know what the minimum abundance can be as a result of thermal freezeout. Because it has a conserved charge, it is also possible to have a larger abundance through generation of an asymmetry. We leave aside this possibility and consider here the abundance of the symmetric component, assuming at first that the \( S \) hyperquark has only HC interactions.

IV.1. \( SU(2)_L \) singlet hyperbaryons

The relic HB abundance is sensitive to the ratio \( \Lambda_{HC}/m_S \), the HC confinement scale over the neutral HQ mass. If \( m_S > \Lambda_{HC} \), there is depletion of the initial \( S \) density through annihilations before confinement, whereas if \( m_S < \Lambda_{HC} \), the \( S \) hyperquarks have a thermal abundance at the confinement phase transition. Once confinement begins, a given \( S \) has a roughly equal probability of forming a hypercolor flux string with a neighboring \( S \) or \( \bar{S} \), leading to roughly equal numbers of hypermesons (that quickly decay away) and HBs. Following the confinement phase transition, there can be further depletion of the HBs by their annihilation.

We estimate the abundance of HBs by solving the Boltzmann equation in the different regimes of temperature described above. The annihilation cross sections for \( S\bar{S} \rightarrow GG \) (annihilation of HQs into hypergluons) and of HBs with their antiparticles are needed. Using refs. \[47, 48\], we find that the first one is

\[
\langle \sigma v \rangle_{S\bar{S} \rightarrow GG} = \frac{\pi \alpha_{HC}^2(m_S)}{4 m_S^2 N_{HC}^3} (N_{HC}^2 - 1)(N_{HC}^2 - 2) \tag{26}
\]

For the gauge coupling we take the four-loop approximation of ref. \[29\] with \( n_f = 0 \) flavors, since we are interested in running only up to the scale \( m_S \), presumed to be the lightest HQ mass in the theory.

For the annihilation of HBs, it is difficult to estimate the cross section due to the strong dynamics and the fact that the HBs are composite states. One possibility is to use the geometric cross section

\[
\langle \sigma v \rangle_{geo} = \frac{\pi}{\mu_*^2} \tag{27}
\]

where \( \mu_* \) is the inverse Bohr radius of the HB, estimated along the lines of ref. \[50\].\(^8\) To extend their method to the regime of strong coupling, we add a linear confining potential \( \sum_{i<j} C_{ij} A_{ij}^2 \) to the Coulomb-like term, to obtain the expectation value of the Hamiltonian for a hydrogen-like wave function \( \psi \sim e^{-\mu_* r^2/2} \) as

\[
\langle H \rangle = \frac{\mu_*^2}{8 m_S} - \frac{5(N_{HC} - N_{HC}^{-1})\alpha_{HC}(\mu_*)}{64} \mu_* + \frac{35 c (N_{HC} - 1) A_{HC}^2}{16 \mu_*} \tag{28}
\]

where the value \( c = 1.9 \) is inferred from refs. \[51, 53\] for the case of \( N_{HC} = 3 \). Minimizing \[28\] with respect to \( \mu_* \) gives an implicit equation for \( \mu_* \) that can be solved by

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\(^7\) We thank Grace Dupuis for computing this using MadGraph.

\(^8\) We disagree with their numerical coefficient for the expectation value of \( 1/r_{ij} \) for the Coulomb-like contribution to the potential.
iteration. The resulting $\mu_s$ as a function of $\Lambda_{nc}$ is shown for several values of $m_S$ in fig. 4. This procedure breaks down when $\Lambda \gtrsim m_S$ because the gauge coupling becomes nonperturbative and the middle term in (25) diverges to large negative values as $\mu_s \rightarrow \Lambda_{nc}$ from above. The approximate values of $\Lambda_{nc}$ where this starts to occur are indicated by heavy dots in fig. 4. Since the dependence of $\log_{10} \mu_s$ on $\log_{10} \Lambda_{nc}$ is very linear (corresponding to the power law $\mu_s \sim \Lambda_{nc}^{0.63}$) in the regions below the dots, we use linear extrapolation to extend our predictions to higher values of $\Lambda_{nc}$. The available final states for HB annihilation almost always include the hypermesons $\pi \pi$, even when they are more quarkonium-like than pion-like (the regime of $m_S \gg \Lambda_{nc}$). The only exception is when $N = 2$ so that mesons and baryons have the same number of constituents.

To compute the relic HB abundance, we numerically solve the Boltzmann equation starting from high temperatures using the $SS \rightarrow GG$ cross section, evolving down to the confinement temperature $T = \Lambda_{nc}$. We assume the confinement transition occurs rapidly and that the initial abundance of HBs at $T = \Lambda_{nc}$ is related to that of HQs by $Y_B = Y_Q/(2N_{nc})$. Taking this as the initial condition for the Boltzmann equation using the HB annihilation cross section, we evolve $Y_B$ to its freeze-out temperature. This approach is generally necessary, rather than the usual analytic approximations, because of the unusual thermal history that often occurs: at the confinement transition, HBs can often be produced starting with a density far exceeding the equilibrium abundance (since the HB mass is $\sim N_{nc}$ times $m_S$). Then the HB annihilations can be in equilibrium, even though a naive treatment would imply that they froze out already at an earlier temperature.

The results are shown in fig. 5 where the fractional abundance of HBs relative to the total observed dark matter is plotted as a function of $\Lambda_{nc}$ for several HQ masses and $N_{nc} = 2, 3, 4$. From conventional thermal freezeout one expects that $\Omega_{HB} \sim 1/(\sigma v)$ as a function of $\Lambda_{nc}$. This explains the simple power-law behavior of $\Omega_{HB}$ in fig. 4(a), since $\mu_s \sim \Lambda_{nc}^{0.63}$ at large $\Lambda_{nc}$. At small $\Lambda_{nc}$ the trend is different because at the confinement temperature $T_c = \Lambda_{nc}$, the initial HB abundance is much higher than the equilibrium abundance. In this situation the conventional dependence is not applicable and we find the different behavior $\Omega_{HB} \sim 1/(\Lambda_{nc} \sigma v)$ in fig. 5.

We find that unless $m_S$ is near the TeV scale (hence not relevant for a diphoton signal at 750 GeV), the symmetric HB component can provide only a subdominant contribution to the total dark matter density. Although the observed density is obtained at low $\Lambda_{nc} \sim 3$ MeV, this region is not viable because of the presence of very light and long-lived glueballs that will disrupt big bang nucleosynthesis for $\Lambda_{nc} \lesssim$ several GeV. But at large $\Lambda_{nc}$, it is possible to obtain the observed abundance. For $m_S \approx 3$ TeV and $\Lambda_{nc} \approx 20$ TeV for example, we find that HBs can constitute all the dark matter, even with no asymmetry. We expect the geometric cross section to provide a lower limit on the true annihilation cross section, which might require a dedicated lattice study to determine with greater certainty. Hence the actual abundance might be smaller than our estimate, although we expect the qualitative dependences to hold. Even a highly subdominant component of HB dark matter could still lead to observable consequences if the HQs have standard model weak interactions, as we describe next.

IV.2. $SU(2)_L$ doublet hyperquarks

As mentioned in section III.2, one way in which the charged HQ $\Psi$ could decay into the neutral one $S$ is by embedding $\Psi$ and $S$ into a fundamental of $SU(5)$ for GUT. This leads to strong constraints on the relic abundance from direct detection. But the same weak interactions could in principle have an impact on the abundance through the annihilations into SM particles. The possible two-body final states include $ZZ, WW, Zh$ and $ff$, depending on the mass of the HQs or HBs. More details on annihilation cross sections are given in appendix C.

However we find that these extra annihilation channels have a negligible effect on the HB abundance, being much weaker than the hypercolor interactions; we can therefore infer the densities from fig. 4. Comparing to the direct detection constraints shown in fig. 6 it can be seen that very light HBs made from HQs with mass $m_S \sim 1$ GeV$\gtrsim \Lambda_{nc}$ can be compatible with the constraints, but heavier ones are ruled out by several orders of magnitude.
V. ELECTROMAGNETIC INTERACTIONS OF NEUTRAL RELICS

Even if the neutral hyperbaryons have no residual strong or weak interactions with nuclei, as would states of the form \( \Phi \Psi \), they inevitably have electromagnetic interactions through the diagrams shown in fig. 6. Although this turns out to be unimportant for direct detection, for completeness we discuss their effects. We estimate these interactions through the diagrams shown in fig. 6. Although not favored by the compatibility of \( \sqrt{s} = 8 \) TeV versus 13 TeV LHC data, a number of authors have shown that the diphoton signal in the 13 TeV data can be accommodated by purely electromagnetic production through photon fusion. The charged hyperquark \( \Psi \) is then neutral under SU(3). If HB number is conserved, then it is compulsory to have neutral hypercolored particles to prevent charged stable hyperbaryons; so we will assume the model is extended with a neutral spinor \( \Phi \) and a scalar \( \Phi \) as previously. Then \( \Psi \) can decay by \( \Psi \rightarrow S f_i f_j \) through interactions of the form

\[
\lambda S^A \Phi \Psi_A + \lambda_{ij} f_i \Phi f_j
\]

where \( f_{i,j} \) are standard model fermions. There are two ways of choosing combinations of SM fermions consistent with gauge invariance, listed in table 1. They correspond to electric charges of 1 or 2 for the mediator and hence of the \( \Psi \), and they imply that \( \Phi \) carries lepton number 2. We can consistently assign the same lepton number to \( S \) so that overall lepton number is conserved by the new interactions. Depending upon the generational structure
TABLE II. Possible combinations of SM fermions coupling to the bosonic mediator $\Phi$, and their charges, for models with uncolored hyperquarks. $L$ and $l$ stand for SU(2)$_L$ doublet and singlet leptons, respectively.

| $f_i$ | $c_{ab}L_a L_b$ | $f_{LR}$ |
|------|----------------|--------|
| $q$  | 1              | 2      |

of the couplings $\lambda_{ij}$ however, there could be violations of individual flavor conservation, or of lepton flavor universality. For example the considerations of section II.3 give $\sqrt{|\lambda_{\mu i} \lambda_{\mu e}|} < 0.6$, for the same choices of mass spectrum, weaker by the factor of $(3N_{hc})^{1/2}$ for the lack of color/hypercolor in the loop.

These models are similar to the favored ones discussed in section III.3 in terms of constraints on the relic HB: they are cosmologically safe, with subdominant relic densities as predicted by figure 5 and unimportant interactions with normal matter through their small loop-induced magnetic moments. Since the scalar mediator couples to lepton pairs instead of being a leptoquark, pairs of charged mesons $\Psi S$ decaying to monoleptons (in the case of the $LL^*$ coupling where one of the particles is a neutrino) or same-sign dileptons would be a collider signature at partonic center of mass energies below 1.5 TeV.

The single-lepton signal is constrained by ATLAS and CMS searches for events with one lepton and missing transverse energy [56, 57]. The more recent ATLAS result limits $m_S \gtrsim 4$ TeV if the couplings $\lambda_{\mu i}$ or $\lambda_{\mu i}$ are of the same order as the SU(2)$_L g$ gauge coupling. The dilepton channel is relatively unconstrained, since ATLAS and CMS searches for same-sign dileptons have so far also required the presence of jets.

VII. CONCLUSIONS

An additional SU($N_{hc}$) gauge group factor is a plausible and economical extension of the standard model. If new matter fields transform in the fundamental of SU($N_{hc}$), the analog of baryon number in the new sector is an issue: if it is conserved then the properties of relic particles must be considered, while if it is broken through renormalizable interactions, interesting constraints can arise from LHC searches for decays associated with these new interactions.

Our considerations are most relevant for models similar to those that can explain the tentative LHC diphoton excess, where we assumed that a charged hyperquark in the fundamental of SU($N_{hc}$) also carries QCD color. Such a particle must decay into standard model states and possibly a neutral state that allows for hyperbaryon number to be conserved, and which is a dark matter candidate. In the case that HB is not conserved, we identified a limited range of renormalizable models, that are summarized in section III.3. Even if the 750 GeV diphoton excess at LHC is only a statistical fluctuation, models consistent with it provide a benchmark for what could be close to the current sensitivity of ATLAS and CMS.

If HB is conserved, the renormalizable models consistent with direct detection constraints and normal baryon conservation are also limited, and turn out to be significantly constrained by LHC leptoquark searches. The viable models have a charged hyperquark $\Psi$ and a scalar mediator $\Phi$ with quantum numbers $(3,2,7/6)$ or $(3,2,1/6)$ under SU(3)$_c \times SU(2)_L \times U(1)_Y$, while the neutral hyperquark $S$ is a singlet. An interesting feature of these models is the presence of composite $\bar{\Psi} S$ leptoquarks (in addition to the heavy fundamental scalar leptoquark $\Phi$), that must have masses $\gtrsim 1$ TeV to satisfy LHC constraints.

An aspect of our work that transcends diphoton signals is the more general possibility that dark matter is a baryon-like state of a new confining sector. The relic density computation for the symmetric component is complicated by the first-order confinement phase transition of the SU($N_{hc}$) sector. We find that hyperquark masses and confinement scales below a TeV, the density is generally much smaller than the observed dark matter density, but for $m_S \sim 3$ TeV, $\Lambda_{hc} \sim 20$ TeV, it could account for all of the dark matter. Searches for anomalous isotopes strongly disfavor the $S$ hyperquark from being colored under QCD, and if it is part of an SU(2)$_L$ doublet, direct dark matter searches limit its abundance to be $\lesssim 10^{-8}$ of the observed DM density, depending upon the hyperbaryon mass. Otherwise the interactions of hyperbaryons with nuclei arise only through loops and give very weak constraints on the model parameters. Such models would be probed more directly through the collider constraints as discussed above.

Note added: as we were completing this work, ref. [60] appeared, which treats astrophysical constraints on models similar to those we have considered, but in the case where the charged HQ $\Psi$ is stable and binds with ordinary quarks to make a neutral relic. According to our analysis in section III.3, such relics are disfavored by anomalous isotope searches, which were not considered in ref. [60].

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The one-loop diagram can be computed exactly, giving

\[ f = 2 + \frac{1}{y^2} x(x+y) \log x^2 \]

\[ - \frac{2[(x+y)^2 - 1][1 - x^2 + xy]}{y^2 x} \ln \frac{2x}{x^2 - y^2 + 1 - r}, \]

where \( f \) is the loop function,

\[ r = \sqrt{[(x-y)^2 - 1][[(x+y)^2 - 1]} \quad (B2) \]

with \( q_\Phi \) the electric charge of \( \Phi \), \( x = m_\Psi/m_\Phi \) and, \( y = m_\Sigma/m_\Phi. \) We plot \( f \) in fig. 8.

To estimate the three-loop diagram, we start by integrating out the \( \Psi \) hyperquark to obtain an Euler-Heisenberg-like effective Lagrangian for the photon-hypergluons vertex \[ (B3) \]

\[ \mathcal{L} = \frac{e g_H^3 d_{abc} }{180 m_\Phi^4 (4\pi)^2} \left[ 14\text{tr}(FG^a G^b G^c) - 5\text{tr}(FG^a)\text{tr}(G^b G^c) \right] \]

where \( F_{\mu\nu}, G^a_{\mu\nu} \) are field strength tensors for photon and hypergluon respectively, \( d_{abc} = 2\text{tr}(\{a, b\}c) \) is the totally symmetric structure constant for \( SU(N_{\text{hc}}) \), and the traces are taken with respect to the Lorentz indices. Then the color factor for the dipole moment diagram is \( d_{abc} \text{tr}(t^at^bt^c)/N_{\text{hc}} = 10/9 \) if \( N_{\text{hc}} = 3 \). We roughly estimate the effect of the Lorentz structure and the additional two loops as giving a factor of \( 9/(16\pi^2)^2 \) to the magnetic moment,

\[ \mu_B \sim \frac{e g_H^3}{180 m_\Phi^4 (4\pi)^2} \cdot \frac{10}{9} \cdot \frac{9}{(16\pi^2)^2} = \frac{e \alpha_H^2}{1152 \pi^3 m_\Phi} \quad (B4) \]

Since the limits from direct detection on this operator are very weak, it is unlikely that a more accurate computation would change our conclusions.

### Appendix B: Dipole moments

In this appendix we estimate the loop-induced interactions of neutral hyperbaryons with photons, relevant for direct detection. The interactions are shown in fig. 6. The one-loop diagram can be computed exactly, giving

\[ \mu_B = \frac{q_\Phi e|\lambda|^2 f}{32\pi^2 m_\Phi} \quad (B1) \]

where \( f \) is the loop function,

\[ f = 2 + \frac{1}{y^2} x(x+y) \log x^2 \]

\[ - \frac{2[(x+y)^2 - 1][1 - x^2 + xy]}{y^2 x} \ln \frac{2x}{x^2 - y^2 + 1 - r}, \]

with \( q_\Phi \) the electric charge of \( \Phi \), \( x = m_\Psi/m_\Phi \) and, \( y = m_\Sigma/m_\Phi. \) We plot \( f \) in fig. 8.

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where \( F_{\mu\nu}, G^a_{\mu\nu} \) are field strength tensors for photon and hypergluon respectively, \( d_{abc} = 2\text{tr}(\{a, b\}c) \) is the totally symmetric structure constant for \( SU(N_{\text{hc}}) \), and the traces are taken with respect to the Lorentz indices. Then the color factor for the dipole moment diagram is \( d_{abc} \text{tr}(t^at^bt^c)/N_{\text{hc}} = 10/9 \) if \( N_{\text{hc}} = 3 \). We roughly estimate the effect of the Lorentz structure and the additional two loops as giving a factor of \( 9/(16\pi^2)^2 \) to the magnetic moment,

\[ \mu_B \sim \frac{e g_H^3}{180 m_\Phi^4 (4\pi)^2} \cdot \frac{10}{9} \cdot \frac{9}{(16\pi^2)^2} = \frac{e \alpha_H^2}{1152 \pi^3 m_\Phi} \quad (B4) \]

Since the limits from direct detection on this operator are very weak, it is unlikely that a more accurate computation would change our conclusions.
here. The relevant s-wave cross sections are

$$\langle \sigma v \rangle_{zz} = \frac{g^4}{64\pi c_W^4 m_S^2} \frac{(1 - x_w^2)^2}{2 - x_w^2}$$

$$\langle \sigma v \rangle_{ww} = \frac{g^4}{64\pi m_S^2} (1 - x_w^2)^2 \left[ \frac{4(4 + 20x_w^2 + 3x_w^4)}{4 - x_w^2} \right] - 4(4x_w + 10x_w^2)(1 + x_w) + 4x_w^4 \left( \frac{1 + x_w - x_h}{x_h - x_w} \right)$$

$$\langle \sigma v \rangle_{zh} = \frac{g^4[(4 - x_h^2)^2 + 2x_w^2(20 - x_h^2) + x_h^4]}{4096\pi c_W^4 m_S^2(4 - x_w^2)^2} \times \sqrt{(4 - x_h^2)^2 - 2x_w^2(4 + x_h^2) + x_h^4}$$

$$\langle \sigma v \rangle_{ff} = \frac{g^4(1 - x_d^2)^2}{8\pi c_W^4 m_S^2(4 - x_w^2)^2} \times [y_{v,f}(2 + x_d^2) + 2g_{af}^2(1 - x_d^2)]$$

where $c_W \equiv \cos \theta_W, x_i = m_i/m_S (i = Z, W, h, f, \Psi)$ and $g_{v,f} (g_{af})$ is the (axial) vector coupling of the fermion $f$ to $Z$ boson. For hyperquarks in fundamental of $SU(N_{HC})$, there is an extra factor of $1/N_{HC}$ for the above formulas taking account of averaging initial degrees of freedom.

For a hyperbaryon with $N_{HC}$ hyperquarks, we expect that there is a coherent enhancement factor of $N_{HC}$ for each gauge interaction vertex of the HB. Then the s-wave annihilation cross section for HB can be rescaled from eq. (C1) as

$$\langle \sigma v \rangle_{zz} = \frac{4}{(N_{HC} + 1)^2} N_{HC}^4 \langle \sigma v \rangle_{zz}$$

$$\langle \sigma v \rangle_{zh} = \frac{4}{(N_{HC} + 1)^2} N_{HC}^2 \langle \sigma v \rangle_{zh}$$

$$\langle \sigma v \rangle_{ff} = \frac{4}{(N_{HC} + 1)^2} N_{HC}^2 \langle \sigma v \rangle_{ff}$$

with $m_S$ replaced by $m_B$ and $x_i$ by $y_i = m_i/m_B (i = Z, W, h, f, \Psi)$. The factor of $4/(N_{HC} + 1)^2$ corrects for the averaging over spin degrees of freedom of the HB. For the WW final state, the rescaling is more complicated because of interference between the s- and t-channel annihilations,
