Dark matter from unification of color and baryon number

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Dark matter from unification of color and baryon number

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We analyze a recently proposed extension of the Standard Model based on the $SU(4) \times SU(2)_L \times U(1)_X$ gauge group, in which baryon number is interpreted as the fourth color and dark matter emerges as a neutral partner of the ordinary quarks under $SU(4)$. We show that under well-motivated minimal flavor-violating assumptions the particle spectrum contains a heavy dark matter candidate which is dominantly the partner of the right-handed top quark. Assuming a standard cosmology, the correct thermal relic density through freeze-out is obtained for dark matter masses around 2 – 3 TeV. We examine the constraints and future prospects for direct and indirect searches for dark matter. We also briefly discuss the LHC phenomenology, which is rich in top quark signatures, and investigate the prospects for discovery at a 100 TeV hadron collider.

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I. INTRODUCTION

The particle identity of the dark matter (DM) is among the most pressing questions confronting particle physics today. It is clear that DM requires an extension of the Standard Model (SM) and it is likely that an understanding of how DM fits into the context of the SM will offer hints about the underlying structure which gave rise to it.

Among its most mysterious properties of DM is the fact that it is long-lived, either being exactly stable or with a lifetime of order that of the Universe itself. Its observational properties are incompatible with it possessing either of the known exactly conserved gauge charges, and thus one would naively expect it to decay quickly. The fact that it is massive and yet (at least to very good approximation) stable provides an important clue about its identity, suggesting that some kind of symmetry is in operation. It is tempting to postulate a connection with the accidental symmetries of the SM, baryon and lepton number, as these are also thought to explain the surprisingly long lifetime of the proton. Indeed, the past years have seen an increase in interest of models gauging $U(1)$ baryon motivated by the proton's lifetime, with early constructions [1–5] paving the way to more complete models [6]. Already, this idea was found in some constructions to lead naturally to theories including DM candidates [7–10].

Extending this idea, in [11] a model unifying gauged baryon number and color into a single non-Abelian gauge group was constructed. The theory is based on the gauge symmetry $SU(4) \times SU(2)_L \times U(1)_X$, with the SM quarks forming quadruplets together with new colorless "quark partner" fields which obtain vector-like masses after $SU(4)$ breaking. The quark partners of the right-handed up-type quarks are SM gauge singlets and have suitable properties to play the role of DM. If the mass of the lightest of such states is chosen to be a few GeV and the gauge structure is supplemented by additional UV interactions, a picture in which the DM number density is determined by a primordial particle-anti-particle asymmetry connected to the asymmetry in baryon number can be realized.

This asymmetric limit is interesting, but also has a few weak points. The quark partners have their own analogue of Yukawa interactions, and thus generically one would not expect the notion of flavor in both the quark and quark partner sectors to be aligned, opening the door for large contributions to flavor-changing neutral currents mediated by the GeV mass DM particle. The need for at least one of the quark partners to have a GeV scale mass precludes the invocation of symmetry-based arguments such as minimal flavor violation (MFV) [12] as a remedy. Perhaps even more unwieldy is the need to introduce an additional sector of light states into which the DM can annihilate (or live with extreme tuning of parameters) to deplete its primordial symmetric component, a generic issue for models of asymmetric DM [13].

These concerns are largely ameliorated if the DM is much heavier and its density is symmetric, resulting from its interactions with the SM quarks freezing out at much higher temperatures. While no longer trying to motivate the observed correspondence between the observed densities of DM and baryons in the Universe, such a limit arises naturally when baryon number and color unify, without the need for ad hoc assumptions or additional ingredients. In the current work, we abandon the connection to the baryon asymmetry and consider the $SU(4)$ model in the limit where all of the quark partners have masses on the order of the $SU(4)$ breaking scale. The D–$\overline{D}$ mixing constraints (derived in [14] for structurally similar leptoquark models) suggest that even in this limit flavor is generically a problem unless there is sufficient alignment between the quark and quark partner Yukawa interactions. As detailed below, we invoke MFV, which results in sufficient alignment for the the first two generations such that D–$\overline{D}$ constraints allow for quark partners with TeV scale masses. The result is a variant of models where the DM is "top-flavored" [15–24], leading to interesting and distinct phenomenology.

This paper is organized as follows: In Section II, we review the unification of baryon number with color, focusing on the features most important for DM. In Section III we compute the rates of annihilation of the DM candidate, as well as its scattering with heavy nuclei, allowing us to identify the regions in which we expect the correct relic density (assuming a standard cosmology) and constraints from the null searches for direct detection of DM. Section IV is devoted to a brief review of the associated collider signals. We summarize our results in Section V.

II. UNIFICATION OF COLOR AND BARYON NUMBER

In this section we provide a brief summary of the most important features of the model for DM and its interactions (for more details see [11]). The underlying gauge structure is:

$$SU(4) \times SU(2)_L \times U(1)_X \ . \tag{1}$$

The SM quarks are promoted to $SU(4)$ quadruplets: $\hat{Q}_L$, $\hat{u}_R$, and $\hat{d}_R$, consisting of the ordinary quark triplets: Q_L, u_R, and d_R, and additional uncolored $SU(4)$ partner fields: $\tilde{Q}_L$, $\tilde{u}_R$ and $\tilde{d}_R$. A generational index should be understood as implicit.

A phenomenologically viable, anomaly-free set of fields is given by:

$$\begin{aligned}
\hat{Q}_L &= (4,2,0)\,, & \hat{u}_R &= \left(4,1,\tfrac{1}{2}\right), & \hat{d}_R &= \left(4,1,-\tfrac{1}{2}\right),\\
Q'_R &= \left(1,2,-\tfrac{1}{2}\right), & u'_L &= (1,1,0)\,, & d'_L &= (1,1,-1)\,,\\
l_L &= \left(1,2,-\tfrac{1}{2}\right), & e_R &= (1,1,-1)\,, & &\\
\hat{\Phi} &= \left(4,1,\tfrac{1}{2}\right), & H &= \left(1,2,\tfrac{1}{2}\right), & &
\end{aligned} \tag{2}$$

where the numbers in parenthesis indicate the representations of $SU(4)$, $SU(2)$, and $U(1)_X$, respectively[1].

The scalar sector contains the $SU(4)$ quadruplet $\hat{\Phi}$, whose vacuum expectation value (VEV),

$$\langle\hat{\Phi}\rangle = \tfrac{1}{\sqrt{2}}\left(0 \;\; 0 \;\; 0 \;\; V\right)^T , \tag{4}$$

breaks the gauge symmetry down to the SM. Hypercharge emerges as a combination of the T^{15} generator of $SU(4)$ and $U(1)_X$,

$$Y = X + \sqrt{\frac{2}{3}}\, T^{15} \ . \tag{5}$$

The $SU(2)_L$ doublet Higgs H breaks the electroweak symmetry down to electromagnetism as usual.

The $SU(4)$ breaking results in seven massive gauge bosons which organize themselves into three complex vector fields G'^{α}_{μ} transforming as a color triplet with mass

$$m_{G'} = \tfrac{1}{2}\, g_4\, V\,, \tag{6}$$

mediating interactions between each SM quark and its partner:

$$\frac{g_4}{\sqrt{2}}\left\{\bar{Q}^{\alpha}_L \not{G}'^{\alpha}\tilde{Q}_L + \bar{u}^{\alpha}_R \not{G}'^{\alpha}\tilde{u}_R + \bar{d}^{\alpha}_R \not{G}'^{\alpha}\tilde{d}_R\right\} + \text{h.c.} \ ; \tag{7}$$

and a neutral Z' gauge boson with mass

$$m_{Z'} = \tfrac{1}{2}\sqrt{g_X^2 + \tfrac{3}{2}g_4^2}\;\; V \ , \tag{8}$$

which couples to pairs of quarks, quark partners or leptons with strength:

$$-\frac{g_Y}{\tan\theta_4}\left[\sqrt{\frac{2}{3}}T^{15} - X\tan^2\theta_4\right]. \tag{9}$$

The angle in $\sin\theta_4 \equiv g_X/\sqrt{g_X^2 + 3g_4^2/2}$ can be determined based on the hypercharge and strong couplings at scale V. For $V \sim$ TeV it is predicted that $\sin\theta_4 \approx 0.28$.

A. Quark Partner Masses and MFV

The masses of the quark partners receive contributions from the VEVs of both $\hat{\Phi}$ and H via Yukawa interactions:

$$\begin{aligned}
& Y_Q^{ab}\,\bar{\hat{Q}}^a_L\,\hat{\Phi}\,Q'^b_R + Y_u^{ab}\,\bar{\hat{u}}^a_R\,\hat{\Phi}\,u'^b_L + \; Y_d^{ab}\,\bar{\hat{d}}^a_R\,\hat{\Phi}\,d'^b_L\\
& + y_u^{ab}\,\bar{\hat{Q}}^a_L\,\tilde{H}\,\hat{u}^b_R + y_d^{ab}\,\bar{\hat{Q}}^a_L\,H\,\hat{d}^b_R\\
& + y'^{ab}_u\,\bar{Q}'^a_R\,\tilde{H}\,u'^b_L + y'^{ab}_d\,\bar{Q}'^a_R\,H\,d'^b_L + \text{h.c.}\ ,
\end{aligned} \tag{10}$$

where the Y couplings marry the quark partners to the spectator fields Q', u', and d'; the y couplings contain the SM Yukawa interactions for the quarks, and the y' couplings lead to mixing between the quark partner singlets and doublets. The result, denoting $\tilde{Q} = (\tilde{U}, \tilde{D})$ and so on for Q', is a pair of 6×6 matrices,

$$\begin{aligned}
&\frac{1}{\sqrt{2}}\left(\bar{\tilde{U}}_L \;\; \bar{u}'_L\right)\begin{pmatrix} Y_Q V & y_u v\\ (y'_u v)^\dagger & (Y_u V)^\dagger\end{pmatrix}\begin{pmatrix} U'_R\\ \tilde{u}_R\end{pmatrix}\\
&+\frac{1}{\sqrt{2}}\left(\bar{\tilde{D}}_L \;\; \bar{d}'_L\right)\begin{pmatrix} Y_Q V & y_d\, v\\ (y'_d v)^\dagger & (Y_d V)^\dagger\end{pmatrix}\begin{pmatrix} D'_R\\ \tilde{d}_R\end{pmatrix} + \text{h.c.}\ ,
\end{aligned} \tag{11}$$

where $v \simeq 246$ GeV is the SM Higgs VEV. The eigenvalues of those two matrices yield the masses of six electrically neutral states (combinations of $\tilde{u}$ and $\tilde{U}$) and six

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[1] We note in passing that it is simple to extend the gauge symmetry to gauge also lepton number: $SU(4) \times SU(2)_L \times U(1)_X \times U(1)_L$. The new anomalies are cancelled by three families of right-handed neutrinos:

$$\begin{aligned}
\hat{Q}_L &= (4,2,0,0)\,, & \hat{u}_R &= \left(4,1,\tfrac{1}{2},0\right), & \hat{d}_R &= \left(4,1,-\tfrac{1}{2},0\right),\\
Q'_R &= \left(1,2,-\tfrac{1}{2},1\right), & u'_L &= (1,1,0,1)\,, & d'_L &= (1,1,-1,1)\,,\\
l_L &= \left(1,2,-\tfrac{1}{2},1\right), & e_R &= (1,1,-1,1)\,, & \nu_R &= (1,1,0,1)\,,\\
S_L &= (1,1,0,-2)\,, & \hat{\Phi} &= \left(4,1,\tfrac{1}{2},-1\right), & H &= \left(1,2,\tfrac{1}{2},0\right),
\end{aligned} \tag{3}$$

where S_L is the additional Higgs needed to break $U(1)_L$. This content allows for new Yukawa terms producing Dirac masses for the neutrinos and a Majorana mass term arising from $S_L\nu_R\nu_R$, thus accommodating a type I seesaw mechanism for neutrino masses.

electric charge minus one states (combinations of $\tilde{d}$ and $\tilde{D}$).

Under the SM flavor symmetries, $\tilde{Q}$, $\tilde{u}$, and $\tilde{d}$ each transform as triplets of $SU(3)_Q$, $SU(3)_u$, and $SU(3)_d$, respectively. The simplest choice[2] is to assign the spectator fields Q', u', and d' to also transform as triplets under $SU(3)_Q$, $SU(3)_u$, and $SU(3)_d$, respectively. MFV then dictates that, to leading order in the spurions y_u and y_d, the remaining Yukawa interactions are given by,

$$Y_Q^{ab} = Y_Q \, \mathbb{1}, \quad Y_u^{ab} = Y_u \, \mathbb{1}, \quad Y_d^{ab} = Y_d \, \mathbb{1} \; , \tag{12}$$

where $\mathbb{1}$ denotes the 3×3 unit matrix, and,

$$y'_u = \eta \; y_u \; , \qquad y'_d = \eta' \; y_d \; . \tag{13}$$

After imposing MFV, the masses of the quark partners are determined by the five parameters: Y_Q, Y_u, Y_d, η, and η', in terms of the SM flavor structure encoded in y_u and y_d.

In the SM quark mass basis, the mass matrices for the partners take the block form,

$$\begin{pmatrix} \bar{\tilde{U}}_L & \bar{u}'_L \end{pmatrix} \begin{pmatrix} M \, \mathbb{1} & m_u \\ \eta \, m_u & m \, \mathbb{1} \end{pmatrix} \begin{pmatrix} U'_R \\ \tilde{u}_R \end{pmatrix} + \begin{pmatrix} \bar{\tilde{D}}_L & \bar{d}'_L \end{pmatrix} \begin{pmatrix} M \, \mathbb{1} & m_d \\ \eta' \, m_d & m' \, \mathbb{1} \end{pmatrix} \begin{pmatrix} D'_R \\ \tilde{d}_R \end{pmatrix} + \text{h.c.} \; , \tag{14}$$

where m_u and m_d are diagonal 3×3 matrices whose entries are the up-type and down-type SM quark masses, whereas $M \equiv Y_Q V/\sqrt{2}$, $m \equiv Y_u V/\sqrt{2}$ and $m' \equiv Y_d V/\sqrt{2}$.

To good approximation (assuming η' is not extremely large) the partners of the first and second generation quarks consist of two degenerate $SU(2)$ doublets of mass $\simeq M$ (along with the partner of the left-handed bottom quark), three degenerate charge -1 singlet states of mass m', and two degenerate neutral singlet states of mass m, with tiny intergenerational mixing and thus negligible contributions to K–$\overline{K}$, B–$\overline{B}$, and D–$\overline{D}$ mixing.

The large top mass results in non-negligible mixing between the $SU(2)$ singlet and doublet top partners, so that their masses are split from M and m. The lighter of the two states, which we denote as χ, is stable due to a global $U(1)$ symmetry left over after the $SU(4)$ breaking and plays the role of DM. Its couplings to the W and Z bosons are controlled by the admixture of the $SU(2)$ doublet, which in turn is controlled by M, m, and η. The mass and gauge eigenstates are related by two mixing angles,

$$\chi_L = \cos\theta_L \; t'_L + \sin\theta_L \; \tilde{T}_L \; , \tag{15}$$

$$\chi_R = \cos\theta_R \; \tilde{t}_R + \sin\theta_R \; T'_R \; . \tag{16}$$

[2] An alternate choice leads to $Y_Q \propto y_u$, $Y_d \propto y_d^\dagger$, and $Y_u, y'_u, y'_d \propto y_u^\dagger$, which would result in large hierarchies in the quark partner masses and require couplings which are nonperturbative.

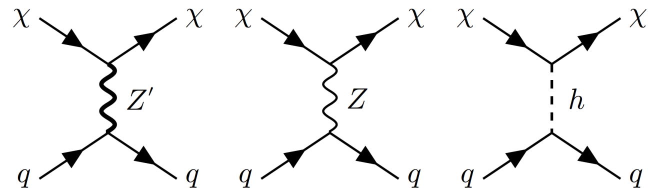

FIG. 1: Representative Feynman diagrams contributing to χ interacting with quarks.

In the limit $M \gg m, m_t$,

$$m_\chi \; \simeq \; m - \frac{m_t^2}{M} \eta \; , \tag{17}$$

$$\sin\theta_R \simeq -\frac{m_t}{M} \eta \; , \qquad \sin\theta_L \simeq -\frac{m_t}{M} \; . \tag{18}$$

As shown below, to evade strong constraints from searches for DM scattering with nuclei, the singlet component should be dominant, i.e. $\theta_R, \theta_L \ll 1$. To simplify our parameter space, we consider $\eta = 1$, for which the two mixing angles are the same. We parameterize the degree of the $SU(2)$ doublet inside χ by

$$\epsilon \equiv \sin\theta_R = \sin\theta_L \ll 1 \; . \tag{19}$$

III. SCATTERING AND ANNIHILATION

In this section we estimate the cross sections for χ to scatter with heavy nuclei or annihilate into SM states.

A. Direct Detection

The DM particle χ interacts with quarks either through the exchange of a heavy Z' or (via its doublet component admixture) the electroweak Z and the SM Higgs boson (see Fig. 1). In the limit of nonrelativistic χ, the effective Lagrangian relevant for the spin-independent cross section can be written as:

$$\mathcal{L}_{\text{eff}} = \sum_q \left[c_1^{(q)} \left(\bar{\chi} \, \chi \right) \left(\bar{q} \, q \right) + c_2^{(q)} \left(\bar{\chi} \gamma^\mu \chi \right) \left(\bar{q} \gamma_\mu q \right) \right], \tag{20}$$

where the coefficients $c_{1,2}^{(q)}$ are given by:

$$c_1^{(q)} = -\frac{\epsilon}{2}\frac{y_t\, y_q}{m_h^2}$$

$$c_2^{(u,c)} = -\frac{g_Y^2}{48}\left[\frac{1}{m_{Z'}^2}\left(\frac{1+\tan^2\theta_4}{\tan^2\theta_4}\right)(2-3\tan^2\theta_4) - \frac{\epsilon^2}{m_Z^2}\left(\frac{2}{\sin^2\theta_W}\right)(3-8\sin^2\theta_W)\right],$$

$$c_2^{(d,s,b)} = -\frac{g_Y^2}{48}\left[\frac{1}{m_{Z'}^2}\left(\frac{1+\tan^2\theta_4}{\tan^2\theta_4}\right)(2+3\tan^2\theta_4) - \frac{\epsilon^2}{m_Z^2}\left(\frac{2}{\sin^2\theta_W}\right)(-3+4\sin^2\theta_W)\right],$$

$$c_2^{(t)} = -\frac{g_Y^2}{48}\left[\frac{1}{m_{Z'}^2}\left(\frac{1+\tan^2\theta_4}{\tan^2\theta_4}\right)(2-3\tan^2\theta_4) - \frac{\epsilon^2}{m_Z^2}\left(\frac{2}{\sin^2\theta_W}\right)(3-8\sin^2\theta_W)\right] - \frac{g_3^2}{8}\frac{1}{m_{G'}^2}. \quad (21)$$

Equation (20) maps onto effective interactions between χ and the nucleon $N = \{p, n\}$:

$$\mathcal{L}_{\rm eff} = \sum_{N=p,n} C_1^{(N)}(\bar\chi\,\chi)\,(\bar N N) + C_2^{(N)}(\bar\chi\gamma^\mu\chi)\,(\bar N\gamma_\mu N) \quad (22)$$

where (e.g. [25]),

$$C_1^{(N)} = \sum_{q=u,d,s} c_1^{(q)}\frac{m_N}{m_q}f_q^{(N)} + \frac{2}{27}f_G^{(N)}\sum_{q=c,b,t} c_1^{(q)}\frac{m_N}{m_q}\,,$$

$$C_2^{(p)} = 2\,c_2^{(u)} + c_2^{(d)}, \qquad C_2^{(n)} = c_2^{(u)} + 2\,c_2^{(d)}, \quad (23)$$

with the coefficients [26] (see also [27]):

$$f_u^{(p)} = 0.023,\ \ f_u^{(n)} = 0.018,\ \ f_d^{(p)} = 0.033,\ \ f_d^{(n)} = 0.042,$$

$$f_s^{(p)} = f_s^{(n)} = 0.26,\ \ f_G^{(p)} = 0.684,\ \ f_G^{(n)} = 0.68\,. \quad (24)$$

The zero-velocity spin-independent cross section for χ to scatter with a nucleon is thus:

$$\sigma_{\rm SI} = \frac{1}{\pi}\frac{m_\chi^2\, m_N^2}{(m_\chi+m_N)^2}\frac{1}{A^2} \times \left[Z\left(C_1^{(p)}+C_2^{(p)}\right) + (A-Z)\left(C_1^{(n)}+C_2^{(n)}\right)\right]^2. \quad (25)$$

Figure 2 shows this cross section as a function of the DM mass for several values of ϵ, assuming that $m_\chi \approx V/3$ (needed for efficient DM annihilation as discussed in the subsequent subsection). Currently, the most stringent limit on $\sigma_{\rm SI}$ for heavy DM comes from the LUX experiment [28]. For m_χ much larger than the mass of a xenon atom, the limit scales simply as $\propto m_\chi$, reflecting the fact that for constant local DM energy density the number density falls as $\propto 1/m_\chi$. Neglecting the subdominant Higgs contribution, for large masses the LUX limit imposes a constraint on V and ϵ:

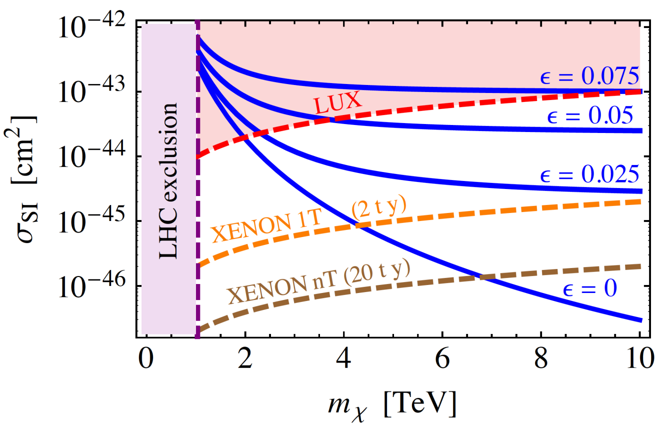

FIG. 2: The spin-independent direct detection cross section as a function of m_χ for various values of ϵ (solid lines) along with current [11, 28] and future [29] exclusion contours (dashed lines). The relation $V \approx 3\,m_\chi$ was adopted to satisfy the DM relic density constraint.

$$\left[\left(\frac{7\ \text{TeV}}{V}\right)^2 + 450\,\epsilon^2\right]^2 \lesssim \frac{m_\chi}{1\ \text{TeV}}\,, \quad (26)$$

generally requiring $V \gtrsim 7$ TeV and $\epsilon \lesssim 0.05$. The bounds on V from LUX are typically stronger than those imposed by null searches for Z' bosons at the LHC [11, 30] or precision electroweak constraints. As shown in Fig. 2, the XENON 1T and XENON nT experiments [29] will probe the available parameter space even further.

B. Annihilation

Due to the severe constraint on the electroweak doublet admixture ϵ in Eq. (26), the DM annihilation into electroweak and Higgs bosons is highly suppressed. The dominant annihilation channels are SM quarks and leptons, mediated by the heavy gauge bosons G' and Z', as shown in Fig. 3. Since the couplings are essentially fixed by the embedding of $SU(3)_c$ and $U(1)_Y$, and the masses are related to one another by the $SU(4)$ breaking scale V, if one assumes a standard thermal history of the Universe the resulting relic density of χ through freeze-out is determined by the values of V and m_χ with little other model dependence.

The resulting relic density is shown in Fig. 4. Cross sections of order $\sim 3\times10^{26}$ cm^3/s are obtained only when m_χ and V are chosen such that annihilation is modestly enhanced by the Z' pole, which happens for $m_\chi \approx V/3$. As is usual in such cases, for fixed V (and thus $m_{Z'}$) there are two values of m_χ for which the thermal relic density is saturated on either side of $m_{Z'}$. Between those two

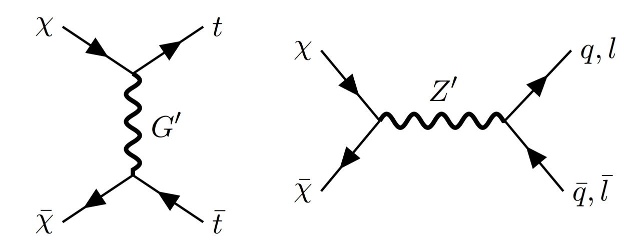

FIG. 3: Representative Feynman diagrams for $\chi\bar{\chi}$ annihilation into SM quarks and leptons.

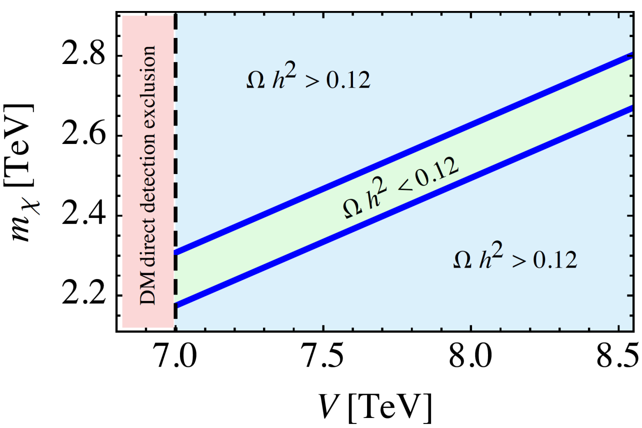

FIG. 4: The relic density of χ relative to the Planck preferred value [31] for a standard cosmological history, with the blue lines saturating $\Omega h^2 \simeq 0.12$, in the plane of V and m_χ.

values of m_χ the cross section is larger and the relic density is typically too small, whereas outside of this range the cross section is too small, and the relic density is in general too large. Limits from indirect detection for this mass range are typically too weak by a few orders of magnitude to provide useful constraints on this parameter space [32, 33].

IV. COLLIDER PHENOMENOLOGY

Given the flavorful nature of the DM, signatures at high energy colliders are typically rich in top quarks. Because of its relatively large coupling to SM quarks and leptons, prospects to observe the Z' at run II of the LHC are good [11], though connecting it to a theory of DM will be more of a challenge. Even for a standard thermal relic, the DM could be heavy enough that there will be no on-shell Z' decays into it, and even if there are $Z' \to \chi\,\bar{\chi}$ decays open, identifying them is likely to prove challenging. It may fall to future high energy colliders to establish the connection between the Z' and DM.

At a future 100 TeV hadron collider, one of the most relevant signatures is pair production of the colored G' from either a gg or $q\bar{q}$ partonic initial state (Fig. 5). The G' will decay into a quark plus a quark partner, resulting in signatures with jets and missing transverse momentum. In particular, decay into χ plus a top quark results in $t+\text{MET}$, or cascading through one of the other quark partners results in $t+j+\text{MET}$. For pair production, the signatures are:

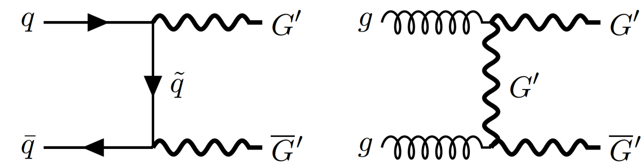

FIG. 5: Representative Feynman diagrams for pair production of $G'\overline{G}'$ at a hadron collider.

$$\text{(a)} \quad p\,p \to G'\,\overline{G}' \to t\,\bar{t} + \text{MET}\ ,$$

$$\text{(b)} \quad p\,p \to G'\,\overline{G}' \to 2\,j + t\,\bar{t} + \text{MET}\ ,$$

where j denotes a light (unflavored or b jet). These signatures are similar to the scalar top or gluino ones from supersymmetry, with mild differences caused by the different spins of the produced particles.

We estimate the rate for pair production of G' by implementing its couplings into a FeynRules model [34, 35], which is used by MadGraph [36] to compute the inclusive cross section. The gg initiated process has a rate determined entirely by gauge invariance under $SU(3)_c$, and thus the only model dependence is the mass of the G' itself. The quark initiated process proceeds via exchange of the quark partners, and thus is sensitive to their mass spectrum. We fix this spectrum by choosing $m = m'$ such that the DM is a canonical thermal relic (see Fig. 4) and $\epsilon \ll 1$ (which requires M to be large enough such that the left-handed quark partners are largely irrelevant).

The resulting cross section is shown in Fig. 6 as a function of the G' mass. For the quark partner masses used to generate this plot, the branching ratio for $G' \to \chi + t$ will be about 1/6, whereas that for $G' \to \chi + j + t$ is around 5/6. The backgrounds for signals such as these at a 100 TeV collider are estimated to be on the order of a femtobarn [37, 38] (after cuts for which the signal events should pass with reasonable efficiency), indicating that G' masses on the order of ~ 7.5 TeV can be probed by this facility. Observation of the G' bosons would be the real clue as to the underlying $SU(4)$ gauge symmetry and its connection to DM.

V. CONCLUSIONS

We have analyzed a novel extension of the Standard Model in which color is unified with baryon number into a single $SU(4)$ gauge group. The theory contains the minimal number of new degrees of freedom consistent with the enlarged gauge symmetry and includes a dark matter

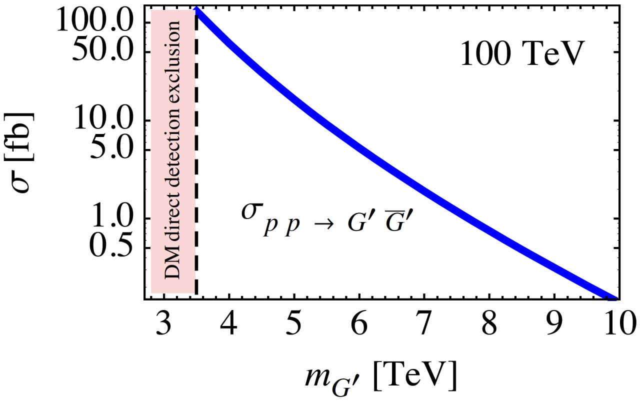

FIG. 6: Cross section for $pp \to G' \overline{G}'$ as a function of the G' gauge boson mass for $E_{\rm CM} = 100$ TeV.

candidate. Constraints from existing searches at colliders and direct detection experiments, along with the dark matter relic abundance, suggest $\sim$ TeV scale masses for these new particles, setting a lower bound on the $SU(4)$ breaking scale of ~ 7 TeV and point to the dark matter being mostly an electroweak singlet.

LHC searches are currently not very constraining, but ultimately have good prospects to detect the Z', which has large coupling to both quarks and leptons. However, its connection to dark matter and the underlying $SU(4)$ symmetry are challenging at the LHC, and would benefit greatly from searches at future colliders, including the 100 TeV pp machine currently under discussion.

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