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Weakly interacting stable hidden sector pions

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An unbroken discrete symmetry, analogous to G parity in QCD, exists in standard model extensions with vectorlike coupling of electroweak $SU(2)$ to "hidden sector" fermions that are confined by a strong gauge force. For an arbitrary irreducible $SU(2)$ representation of the hidden sector fermions, the lightest hidden sector states form an isotriplet of "pions" with calculable mass splittings and couplings to standard model fields. The parity can be extended to fermions in real representations of color $SU(3)$, and can provide dark matter candidates with distinct collider signatures.

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I. INTRODUCTION

Cosmological and astrophysical observations indicate the existence of dark matter. This motivates the investigation of standard model (SM) extensions with new stable particles. For example, in theories with an exact parity symmetry, the lightest parity-odd particle is stable against decay, and can act as dark matter. Here we describe the implications of a parity selection rule that naturally arises in a "hidden sector" involving strongly coupled fermions, e.g. a QCD-like theory with large confinement scale.

At the renormalizable level, such a hidden sector can communicate with the standard model only through gauge interactions: $SU(3)_c$, $SU(2)_W$, or $U(1)_Y$. To avoid spontaneous breaking of these symmetries at a high scale we focus on vectorlike gauging, i.e., identical coupling of the gauge fields to left- and right-handed fermions.

We will see that the class of models with gauge-confined fermions coupled vectorially to $SU(2)_W$ has some remarkably simple universal features. A parity symmetry acting only on the hidden sector fields is left unbroken. The new parity symmetry is an analog of G parity in QCD [1], which is broken in the SM by the presence of $U(1)_Y$ and nonvectorial coupling of $SU(2)_W$. Below the hidden sector confinement scale, the lightest physical states are Nambu-Goldstone bosons (NGBs) of the spontaneously broken chiral symmetries. In the limit of vanishing electroweak coupling ($g_2 \to 0$), the pions are degenerate in mass, as implied by exact global symmetries. At finite coupling, and after electroweak symmetry breaking, radiative corrections induce mass splittings, and the electric-neutral component of an $SU(2)_W$ triplet is the "lightest G-odd particle" (LGP). If such a hidden sector exists, this particle can be a dark matter candidate. Its mass and standard model couplings are calculable in terms of a free parameter representing the confinement scale of the hidden sector fermions.

Models with vectorlike coupling of standard model gauge fields to hidden sector matter have previously been considered from the point of view of dark matter [2] and collider [3] phenomenology. In a large class of models, our conclusions differ from previous work. We demonstrate how stable particles naturally arise in such models without *ad hoc* assumptions, and provide a chiral Lagrangian framework for systematically computing properties of the new particles.

II. FERMION DESCRIPTION

Let us extend the standard model by introducing Dirac fermions in a complex representation of a hidden sector gauge group, e.g. the fundamental of $SU(N)_h$ with $N \geq 3$, i.e., a QCD-like theory. Suppose there are $n_f = 2j+1$ flavors transforming in an irreducible isospin j representation of $SU(2)_W$.[1] The basic Lagrangian is

$$\mathcal{L} = \mathcal{L}_{\rm SM} - \tfrac{1}{4}(\hat{F}^a_{\mu\nu})^2 + \bar{\hat{\psi}}(i\partial\!\!\!/ + \hat{g}\hat{A\!\!\!/}^b \hat{t}^b + g_2 W\!\!\!\!/^a J^a)\hat{\psi}, \quad (1)$$

where hatted quantities refer to the hidden sector. The Hermitian matrices $\hat{t}^b$, $b = 1\cdots N^2-1$ denote generators in the fundamental $SU(N)$ representation and J^a, $a = 1\cdots 3$ are $n_f \times n_f$ isospin matrices.

All representations of $SU(2)$ are real. The matrix S that relates the original and conjugated generators, $S^\dagger J^a S = -J^{a*}$, can be expressed in terms of a group element [4]: $S = \exp[i\pi J^2]$. On the basis of isospin j states, $|jm\rangle$ with $m = j\cdots -j$, we have

$$S_{m_1 m_2} = (-1)^{j+m_2}\delta_{m_1,-m_2} = (\exp[i\pi J^2])_{m_1 m_2}. \quad (2)$$

For example, for $j = 1/2$ we have $J^a = \tau^a/2$ and $S = i\tau^2$, where τ^a are the usual Pauli matrices.[2] We use the reality of $SU(2)$ representations to define a modified charge conjugation operation, "G parity,"

[1]We assume that N and n_f are such that spontaneous chiral symmetry breaking occurs. We are not here concerned with issues such as perturbativity and unification at high scales.

[2]This example was discussed in Ref. [3], where it was incorrectly stated that stable particles are absent for $j > 1/2$.

$$\hat{\psi} \xrightarrow{G} S\hat{\psi}^{\mathcal{C}} = Si\gamma^2\hat{\psi}^*, \quad \hat{A}^b\hat{t}^b \xrightarrow{G} (\hat{A}^b)^{\mathcal{C}}\hat{t}^b = \hat{A}^b(-\hat{t}^{b*}), \quad (3)$$

with $\mathcal{C}$ denoting "ordinary" charge conjugation and S acting on the fermion flavor index. All SM fields, in particular W^a_μ, are left invariant. It is readily verified that the Lagrangian (1) is invariant under the transformation (3).[3] G parity is a good quantum number of the theory, and all SM particles are G even.

III. PION DESCRIPTION

The $SU(N)_h$ gauge theory is assumed to have a behavior similar to QCD. The gauge coupling becomes strong in the infrared, triggering confinement and chiral symmetry breaking at a scale Λ_h. Below Λ_h, the effective theory is described by "pions" which are NGBs associated with the spontaneously broken global flavor symmetry of the hidden sector: $SU(2j+1)_L \times SU(2j+1)_R \times U(1)_B \to SU(2j+1)_V \times U(1)_B$. Here $U(1)_B$ is the unbroken dark baryon number symmetry.

The pions are collected into the $SU(2j+1)$ matrix field $U \equiv e^{2i\tilde{\Pi}/f_\Pi}$, where f_Π is the analog of $f_\pi \approx 93$ MeV in QCD. The $(2j+1)^2 - 1 = \sum_{J=1}^{2j}(2J+1)$ pions can be classified in irreducible representations, J, of $SU(2)_W$. The generators $t^{(JM)}$ associated with pions of definite total isospin J and electric charge M can be expressed using Clebsch-Gordan coefficients for $j \times j$ as

$$t^{(JM)}_{m_1 m_2} = (-1)^{j-m_2}\langle jjJM|jjm_1, -m_2\rangle. \quad (4)$$

The generators (4) are normalized as $\mathrm{Tr}[t^{(JM)}t^{(J'M')}] = (-1)^M\delta^{JJ'}\delta^{M,-M'}$, are real by definition, and satisfy $[t^{(JM)}]^T = (-)^M t^{(J,-M)}$. The canonically normalized isospin generators appearing in (1) are $J^3 = \sqrt{C(j)}t^{(10)}$, $J^\pm/\sqrt{2} = \mp\sqrt{C(j)}t^{(1\pm1)}$, where $J^\pm = J^1 \pm iJ^2$, and $C(j) = j(j+1)(2j+1)/3$ is the normalization constant for the isospin-j representation: $\mathrm{Tr}(J^aJ^b) = C(j)\delta^{ab}$. The pion field can be expanded as

$$\tilde{\Pi} = \sum_{i=1}^{(2j+1)^2-1}\Pi^i t^i = \sum_{J=1}^{2j}\sum_{M=-J}^{J}\Pi^{(JM)}t^{(JM)}. \quad (5)$$

Under the transformation (3) and using (2) we have

$$\Pi^{(JM)} \xrightarrow{G} (-1)^J\Pi^{(JM)}, \quad (6)$$

so that pions with odd (even) J are odd (even) under G parity. Note that G parity generalizes the notion of pion number parity in QCD to arbitrary number of flavors.

[3]The path integral measure is similarly invariant, so that the parity is well defined at the quantum level.

IV. SPECTRUM

The leading, two-derivative, term in the symmetric chiral Lagrangian is

$$\mathcal{L}_2 = \frac{f_\Pi^2}{4}\mathrm{Tr}[D_\mu U D^\mu U^\dagger], \quad (7)$$

where $D_\mu U \equiv \partial_\mu U - ig_2 W^a_\mu[J^a, U]$. This term describes massless, self-interacting scalar fields. Note that the neutral pions $\Pi^{(J0)}$ correspond to generators commuting with J^3, and hence have no direct interactions with the physical photon or Z boson. The gauge coupling g_2 explicitly breaks the global chiral symmetry. At one loop, a quadratically divergent counterterm is required,

$$\mathcal{L}_{\mathrm{g}} \sim g_2^2\Lambda_h^2 f_\Pi^2\,\mathrm{Tr}(J^aUJ^aU^\dagger). \quad (8)$$

This loop effect lifts the degeneracy among multiplets,

$$m^2_{\Pi^{(JQ)}} \sim J(J+1)\alpha_2\Lambda_h^2, \quad (9)$$

where $g_2^2 \equiv 4\pi\alpha_2$. The hidden sector couples to the SM Higgs field via the $SU(2)_W$ gauge field intermediary. After electroweak symmetry breaking, such interactions induce a finite splitting between components of each multiplet. In the limit $m_\Pi \gg m_W$, $m_{\Pi^{(JQ)}} - m_{\Pi^{(J0)}} \approx \alpha_2 Q^2 m_W \sin^2\frac{\theta_W}{2}$ [2,5]. In particular, $m_{\Pi^{(1\pm1)}} - m_{\Pi^{(10)}} \approx 170$ MeV, leaving $\Pi^{(10)}$ as the LGP.

V. INTERACTIONS

The interactions of the pions among themselves, and with SM fields, are constrained by chiral symmetries of the new strong interaction. Taking over results familiar from QCD, the symmetric Lagrangian can be expanded in powers of $1/f_\Pi$: $\mathcal{L} = \mathcal{L}_2 + \mathcal{L}_4 + \cdots$. The leading term is displayed in (7). Of particular interest at four-derivative order is the Wess-Zumino-Witten (WZW) interaction, containing terms with odd pion number,

$$\mathcal{L}_4 = \frac{Ng_2^2}{16\pi^2 f\Pi}\epsilon^{\mu\nu\rho\sigma}\mathrm{Tr}[\tilde{\Pi}W_{\mu\nu}W_{\rho\sigma} + \mathcal{O}(\tilde{\Pi}^3)] + \cdots. \quad (10)$$

Let us consider the decays of the lightest states, focusing for simplicity on the limit $m_\Pi \gg m_W$. The charged members of the ground-state multiplet decay with weak strength into the LGP via QCD pion emission:

$$\Gamma(\Pi^{(1\pm1)} \to \Pi^{(10)} + \pi^\pm) = \frac{4G_F^2}{\pi}\Delta^2\sqrt{\Delta^2 - m_\pi^2}f_\pi^2, \quad (11)$$

where $\Delta = m_{\Pi^{(1\pm1)}} - m_{\Pi^{(10)}}$ and we have assumed $m_\Pi\Delta \gg m_\pi^2$. In this limit the decay length $c\tau \approx 5$ cm is independent of m_Π. The lightest G-even pions, $\Pi^{(2M)}$, decay into electroweak bosons via the WZW term (10). The relevant trace is

$$C(j)\,\mathrm{Tr}[t^{(2M)}\{t^{(1L)},t^{(1L')}\}] = (-)^M\langle 112M|11-L'-L\rangle \frac{4}{\sqrt{15}}[(j-1/2)j(j+1/2) \times (j+1)(j+3/2)]^{1/2}. \quad (12)$$

The amplitude is sensitive both to the number of colors N and to the isospin j of the underlying fermions. The rate scales as $\Gamma(\Pi^{(2M)}) \sim \alpha_2^2 m_\Pi^3/(4\pi)^3 f_\Pi^2$. Note that the process $\Pi^{(1M)} \to WW$ (the analog of $\pi^0 \to \gamma\gamma$ in the SM) is forbidden by the vanishing of the relevant isospin trace (“d symbol”).

The next-to-lightest G-odd pions, $\Pi^{(3M)}$, decay into the ground-state multiplet via loops containing G-even pions, $\Pi^{(2M)}$. The interaction vertices are obtained by expanding (7) and (8). The decay rate scales as $\Gamma(\Pi^{(3M)} \to \Pi^{(1M')} + 2W) \sim \alpha_2^2 m_\Pi^5/(4\pi)^5 f_\Pi^4$. For $J \geq 3$, loops containing G-even pions mediate the decays $\Pi^{(JM)} \to \Pi^{(J-2,M')} + 2W$. The end result is a multi-W final state. For odd J, the stable $\Pi^{(10)}$ remains, possibly after decay of the long-lived $\Pi^{(1\pm1)}$ as above.

VI. COUPLING TO $SU(3)_c$

The parity operation (3) can be extended to real representations of $SU(3)_c$. [Extension to $U(1)_Y$ would require embedding the gauge fields inside a larger, real representation.]

Consider, e.g., hidden sector fermions transforming in the adjoint representation of $SU(3)_c$, and the fundamental representation of $SU(2)_W$. In the usual basis where adjoint generators are purely imaginary, the generalized G parity for this model is defined as in (3), where now $S = \mathbb{1}_8 \otimes i\tau^2$ acting on $SU(3)$ and $SU(2)$ indices. Using the product decomposition $\mathbf{8}\times\mathbf{8} = (\mathbf{8}+\mathbf{10}+\overline{\mathbf{10}})_{\mathrm{A}} + (\mathbf{1}+\mathbf{8}+\mathbf{27})_{\mathrm{S}}$ (“A” and “S” denoting antisymmetric and symmetric tensors) and $\mathbf{2}\times\mathbf{2} = \mathbf{1}+\mathbf{3}$, we can express the $16^2-1 = 255$ pion generators as products of $SU(3)$ and $SU(2)$ generators. The resulting symmetric (antisymmetric) pure $SU(3)$ generators are G even (G odd) and the decomposition into representations of $SU(3)_c \times SU(2)_W$ with the corresponding G parity is $(\mathbf{8},\mathbf{2})\times(\mathbf{8},\mathbf{2}) - (\mathbf{1},\mathbf{1}) = (\mathbf{1},\mathbf{3})^- + (\mathbf{8}_{\mathrm{A}},\mathbf{1})^- + (\mathbf{8}_{\mathrm{S}},\mathbf{1})^+ + (\mathbf{8}_{\mathrm{A}},\mathbf{3})^+ + (\mathbf{8}_{\mathrm{S}},\mathbf{3})^- + \cdots$. The pion masses are proportional to the sum of second-order Casimir invariants as in (9), $m^2_{\Pi^{(d_3,d_2)}} \sim \Lambda_h^2[\alpha_3 C_2(d_3) + \alpha_2 C_2(d_2)]$, with $C_2(3) = 2$ for $SU(2)$, and $C_2(8) = 3$ for $SU(3)$. The lightest states, $(\mathbf{1},\mathbf{3})^-$, are again an $SU(2)_W$ triplet of G-odd pions.

Operators such as (7), (8), and (10) again mediate transitions between pion states, after appropriate substitution to account for $SU(3)_c \times SU(2)_W$ gauging. The decays of the two lightest G-even multiplets $(\mathbf{8}_{\mathrm{S}}, 1)^+$ and $(\mathbf{8}_{\mathrm{A}},\mathbf{3})^+$, proceed primarily through WZW interactions. The two next-to-lightest G-odd pions, $(\mathbf{8}_{\mathrm{A}},\mathbf{1})^-$ and $(\mathbf{8}_{\mathrm{S}},\mathbf{3})^-$, decay to the LGP plus two SM gauge bosons through loop diagrams with G-even pions. The spectrum and interactions in this model can produce interesting collider signatures. For example, the pair production of the next-to-lightest G-odd particle $(\mathbf{8}_{\mathrm{A}},\mathbf{1})^-$ leads to a final state with 2 photons + 2 jets + missing transverse energy [5], with displaced vertices from intermediate metastable electrically charged $(\mathbf{1},\mathbf{3})^-$ states.

Other possibilities for the hidden sector include reducible representations of $SU(2)_W$, i.e., multiple “generations.” These cases might be interesting to investigate from the standpoint of SM gauge-singlet dark matter (interacting via higher-dimension operators), or as nearly degenerate isotriplets. Another possibility (not relevant to weakly coupled dark matter) is to consider $SU(2)_W$ singlet fermions in an irreducible real representation of $SU(3)_c$. Various models can be embedded in five-dimensional gauge theory constructions.

VII. PECCEI-QUINN SYMMETRY AND AXION

We have so far neglected two gauge invariant terms in the Lagrangian (1) that can appear at the renormalizable level. The first is a theta term for the hidden gauge fields, and the second a bare mass term for the fermions,[4]

$$\mathcal{L}_\theta = \theta\epsilon^{\mu\nu\rho\sigma}\hat{F}^a_{\mu\nu}\hat{F}^a_{\rho\sigma}, \qquad \mathcal{L}_m = -m_\psi \bar{\hat{\psi}}\,\hat{\psi}. \quad (13)$$

For nonzero m_ψ and θ, these terms give rise to P and T violation in the hidden sector but do not affect G-parity conservation.[5] For simplicity we focus on the case $m_\psi \ll \Lambda_h$. For $m_\psi \gg \Lambda_h$, the lightest G-odd states become isotriplet nonrelativistic “quarkonium.”

It is natural to consider whether some mechanism suppresses P and T violation in the hidden sector, as happens in QCD. Note that at $m_\psi = 0$, there is a Peccei-Quinn (PQ) symmetry present at the classical level, $\hat{\psi} \to e^{i\alpha\gamma_5}\hat{\psi}$. The PQ symmetry can be preserved at $m_\psi \neq 0$ if the mass is generated by spontaneous symmetry breaking. Consider a scalar field σ that transforms under the PQ symmetry as $\sigma \to e^{-2i\alpha}\sigma$, and with interactions:

$$\mathcal{L}_\sigma = |\partial_\mu\sigma|^2 - V(\sigma) - \lambda\sigma\bar{\hat{\psi}}_L\hat{\psi}_R + \text{H.c.}, \quad (14)$$

where $V(\sigma)$ is such that σ acquires a vacuum expectation value, $\langle\sigma\rangle = f_a$. The mass parameter is then $m_\psi = \lambda f_a$, and with $\sigma(x) \sim f_a e^{ia(x)/\sqrt{2}f_a}$, the low-energy spectrum includes an axion $a(x)$ which is massless at the classical level.

As with the QCD axion, we assume that the $SU(N)_h$ anomaly generates an effective potential for $a(x)$ with minimum such that the effective θ term vanishes. Gauge invariance prevents the hidden axion from direct interaction

[4]We have used the freedom to perform a chiral rotation of the fermions and redefine θ to eliminate any $\bar{\hat{\psi}}\gamma_5\hat{\psi}$ term.

[5]A Majorana mass term violating G parity is forbidden by gauge invariance for complex representations of $SU(N)_h$.

with SM fermions.[6] The physical axion acquires mass and interactions through mixing with the NGB of the $U(1)_A$ global flavor symmetry in the hidden sector,

$$\mathcal{L}_a = \frac{m_a^2}{2f_\Pi^2}[\Pi^i]^2 a^2 - \frac{NC(j)g_2^2}{32\pi^2\sqrt{2}f_a} a\epsilon^{\mu\nu\rho\sigma} W^a_{\mu\nu} W^a_{\rho\sigma} + \cdots, \tag{15}$$

where $m_a^2 \sim m_\psi \Lambda_h f_\Pi^2/f_a^2$. Note that the $SU(2)_W$ couplings to the hidden sector do not break the $U(1)_{PQ}$ symmetry, and so do not contribute to the axion mass: only in the limit $\sqrt{m_\psi \Lambda_h} \gg g_2 \Lambda_h$ does the QCD-like relation $m_a f_a \sim m_\Pi f_\Pi$ hold.

VIII. DISCUSSION

Hidden sector fermions coupled vectorially to SM gauge fields are an ingredient in typical axion models addressing the strong CP problem [6]. It is natural to inquire whether other such hidden sectors exist. We have considered approximately massless Dirac fermions in a complex representation of a QCD-like gauge group, and an arbitrary vectorlike coupling to $SU(2)_W$. The lightest state of such a hidden sector is the neutral component of an approximately degenerate isotriplet pion. In contrast to the C and CP symmetries in the SM, there are no gauge invariant renormalizable operators that violate the new G parity.

Could weakly interacting stable pions (WISPs) be a component of cosmological dark matter, produced thermally in the early universe? A lifetime of order $\tau_{\rm universe} \sim 10^{10}$ yr imposes tight constraints and we must consider whether G parity could be broken either explicitly (by unknown UV physics) or spontaneously (from the choice of vacuum). Suppose that the SM extension (1) is viewed as an effective theory valid up to some scale $\Lambda_{\rm UV}$. Corrections to the renormalizable Lagrangian can appear at dimension five. Two such operators, $B_{\mu\nu}\bar{\hat{\psi}}\sigma^{\mu\nu}\hat{\psi}/\Lambda_{\rm UV}$ and $H^\dagger t^a H \bar{\hat{\psi}} J^a \hat{\psi}/\Lambda_{\rm UV}$, violate G parity. Even for $\Lambda_{\rm UV}$ of the order of the Planck scale, an additional suppression would be necessary to ensure cosmological stability. Note however that these operators violate the PQ symmetry present at the renormalizable level. Enforcing this symmetry implies the appearance of the σ field, leading to an additional suppression $\langle\sigma(x)\rangle/\Lambda_{\rm UV} \sim f_a/\Lambda_{\rm UV}$. Then $\tau_{\rm LGP} \sim \Lambda_{\rm UV}^4/f_\Pi^3 f_a^2 \sim \tau_{\rm universe} \times 10^8 [f_\Pi/f_a]^2$, for $\Lambda_{\rm UV} \sim 10^{16}$ GeV and $f_\Pi \sim 10^3$ GeV. Provided that f_a/f_Π is not too large, the lifetime of the LGP is plausibly long enough to be a thermal dark matter candidate.

Consider the vacuum alignment question: can the fermion condensate $\Sigma_0^{ij} \sim \langle \bar{\hat{\psi}}_L^i \hat{\psi}_R^j \rangle$ spontaneously break G parity? For massless fermions, the effective potential that determines Σ_0 begins at order g_2^2. As indicated by the positive pion masses (9), the potential has a minimum at $\Sigma_0 \propto \mathbb{1}$. This conclusion is unaffected if a bare fermion mass (13) is included in the hidden sector: the gauged vectorlike symmetries are unbroken [7]. Interactions with the SM lead to $\mathcal{O}(g_2^4)$ corrections to the effective potential. Since G parity is not explicitly broken, the potential is invariant under G parity, and by a simple continuity theorem [8], the shift (if any) of Σ_0 is similarly invariant. Barring a conspiracy of higher-order terms, G parity should not be spontaneously broken.

What about hidden sector baryons? These are also stable particles and potentially contribute to the dark matter relic abundance. In the absence of hidden baryon nonconservation, baryon-antibaryon pairs are created in thermal equilibrium. At freeze-out, the energy density of baryons to pions scales as [5] $\Omega_B/\Omega_\Pi \sim \langle\sigma v\rangle_\Pi/\langle\sigma v\rangle_B \sim 4\pi\alpha_2^2 \sim 10^{-2}$, indicating that the baryons play a subdominant, but non-negligible role. (Here σ denotes annihilation cross section.) In cases where N and j forbid a neutral baryon, however, the stable baryons can be fractionally charged, imposing tight constraints. Models of (techni-) baryonic dark matter have been investigated in [9].

What are the existing constraints on WISPs? The impact of a vectorlike hidden sector on low-energy precision measurements is minimal [3,10]. A search for charged fermions decaying to missing energy at LEP constrains $m_{\rm LGP} \gtrsim 88$ GeV [11]. Studies of nearly degenerate fermion isotriplets suggest a potential sensitivity up to $\sim$100 GeV at the Tevatron [12], and somewhat higher mass at the LHC [13]; production channels involving $SU(3)_c$ octets would increase the reach. The precise values of the bounds will depend on production cross sections for isotriplet scalars versus fermions. Studies of $SU(2)_W$ multiplets of simple scalar dark matter particles [2] suggest a bound $m_{\rm LGP} \lesssim 1$–2 TeV in order to not overclose the Universe, and a spin-independent scattering cross section on a nucleon $\sim 10^{-45}$ cm^2; such particles should have so far evaded direct dark matter detection experiments [14]. In all of these constraints, precise values are affected by pion self-interactions, and possibly by axion or color $SU(3)_c$ interactions. We conclude that a large window $\sim$100 GeV–few TeV potentially exists for the LGP mass. A more thorough analysis will be presented elsewhere.

What are the constraints on the hidden sector axion? The lifetime, $\tau_a \sim [\alpha^2 m_a^3/f_a^2]^{-1}$, should be such that axions produced thermally in the early universe decay before big bang nucleosynthesis. This avoids constraints from elemental abundances, CMB distortions, and diffuse photon backgrounds [15]. Assuming for illustration $m_\psi \sim m_\Pi \sim 100$ GeV, this yields $f_a \lesssim 10^{4-5} f_\Pi$. Note that if the axion mass is not small compared to $g_2 f_\Pi$, the LGP can annihilate into two axions with substantial branching fraction, influencing the thermal history and potential dark matter annihilation signals.

[6] The PQ symmetry could be extended to accommodate a right-handed neutrino coupling by assigning a common vector PQ charge to all SM leptons.

In conclusion, we have argued that an unbroken and nontrivial discrete symmetry emerges in a class of models with vectorlike SM gauging of fermions with QCD-like strong dynamics. The mechanism leads to testable and distinct dark matter and LHC phenomenologies.

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