Abstract

We discuss a dark family of lepton-like particles with their own “private” gauge bosons $X_\mu$ and $C_\mu$ under a local $SU'(2) \times U'(1)$ symmetry. The Higgs doublet would couple in the standard way to the left-handed $SU'(2)$ doublet $\Psi_L$ and right-handed singlets $\psi_{1,R}, \psi_{2,R}$, but not to the extra gauge bosons. This reduces the electroweak-type gauge symmetries from $SU'(2) \times U'(1) \times SU(2) \times U(1)$ to the diagonal (vector-like) $SU'(2) \times U(1)$. The “dark leptons” would contribute to the dark matter and interact with Standard Model matter through the Higgs portal. The dark gauge bosons $X_\mu$ and $C_\mu$ remain massless with their energy densities decreasing like $a^{-4}$. This defines a dark matter model with internal interactions. The Standard Model Higgs boson aligns the two electroweak-type symmetry groups in the visible and dark sectors and generates the masses in both sectors. We also identify charge assignments for $SU'(2) \times U'(1)$ in the dark sector which allow for the formation of dark atoms as bound states of dark lepton-like particles. The simplest single-component dark matter version of the model predicts a dark matter mass around 96 GeV, but the corresponding nucleon recoil cross section of $1.2 \times 10^{-44} \text{cm}^2$ is ruled out by the xenon based experiments. However, multi-component models or models with a dark $SU(2)$ doublet mediator instead of the Higgs portal would still be viable.

Keywords: Dark matter, Extensions of the Standard Model

PACS: 95.35.+d, 12.60.-i

1. Introduction

The identification of the dark matter which dominates the large scale structure in the universe remains an open problem. Higgs exchange remains an interesting and observationally viable option for non-gravitational interactions between dark matter and ordinary matter [1, 2, 3, 4, 5, 6, 7, 8]. Higgs portal models through couplings to the scalar product $H^*H$ of the Higgs doublet $H = (H^+, H^0)$ have been discussed extensively for bosonic [9, 10, 11, 12, 13, 14, 15, 16, 17, 18, 19, 20, 21, 22, 23, 24, 25, 26] and fermionic [18, 27, 25, 26] dark matter, and will be tested over a wide mass range within the next few years.

Mass terms in the dark sector are usually inserted by hand or generated by a separate symmetry breaking mechanism. However, it is an intriguing question whether $SU(2)$ breaking by the Standard Model (SM) Higgs could also generate the dark matter masses without violating the Standard Model gauge symmetries, and yet be safe from constraints arising through couplings to the electroweak gauge bosons. We will answer this question affirmatively. The key idea is to have an a priori separate $SU'(2) \times U'(1)$ gauge symmetry in the dark sector with its own gauge bosons $X_\mu$ and $C_\mu$. The standard chiral Yukawa couplings of dark left-handed $SU'(2)$ lepton-like doublets and right-handed $SU'(2)$ singlets with the Higgs doublet then break the gauge symmetry of the coupled SM+dark matter model to $SU_c(3) \times SU_6(2) \times U_Y(1)$, because the Higgs couplings align the local $SU(2) \times U(1)$ gauge transformations in the visible and dark sectors. The “private” dark sector gauge bosons do not couple to the Higgs doublet and therefore remain massless, with their cosmological energy densities after decoupling in the dark sector decaying with the Friedmann-Robertson-Walker (FRW) scale factor $a(t)$ proportionally to $a^{-4}$. On the other hand, the dark lepton-like particles (for simplicity de-
noted as dark leptons in the following) acquire mass terms through their SM-like Yukawa couplings to the Higgs. The dark leptons therefore constitute the dark matter in the model, assuming that leptogenesis also had an analog in the dark sector. Neutrality under the internal gauge interactions of the dark sector then implies that the dark matter will consist of dark atoms. The resulting dark matter model is therefore similar to the dark atom models, see [30, 31, 32, 33] and references there.

The primary new idea is the Higgs alignment between visible and dark gauge groups.

The model is introduced in Sec. 2. The formation of dark particles in (1), but the generalization to more genera-
Here \( q_2 \) and \( q_1 \) are the gauge couplings of the dark \( SU'(2) \times U'(1) \) symmetry group, and the covariant derivatives on the dark matter fields are

\[
D_p \psi_L(x) = \partial_p \psi_L(x) - iq_2 X_p(x) \cdot \sigma \psi_L(x) - i q_1 \frac{1}{2} Y_L'(x) C_p(x) \psi_L(x),
\]

(6)

\[
D_p Y_{L,R}(x) = \left( \partial_p - i \frac{q_1}{2} Y_{L,R}'(x) C_p(x) \right) \psi_L R(x),
\]

(7)

\[
Y_{L,R} = Y_L^* = -y',
\]

(8)

The reader will certainly have noticed that the assumed preference for left-chirality also in the dark sector is not determined by any experimental observation and therefore ambiguous with our current knowledge. Instead, we could just as well use right-handed \( SU'(2) \) doublets and left-handed singlets in the dark sector. This would be compelling from an enhanced symmetry point of view. Opposite chirality in the dark sector could restore \( CP \) and time-reversal symmetry in particle physics through mappings between visible sector and dark sector fields. \( CP \) and time reversal symmetry might then only be hidden from us through the very weak Higgs-coupling between the visible sector and the dark sector. The model would then provide a Higgs portal realization of the old idea of a \( CP \) mirror of the Standard Model \( SU(3 \timesSU)\). However, currently we can constrain the dark sector only through gravitational effects and direct and indirect search limits. We cannot currently distinguish between a \( V-A \) or \( V+A \) preference in the dark sector, and at this stage a \( V-A \) formulation is as good as a \( V+A \) formulation.

There are obviously infinitely many possible ramifications of this model by also including corresponding families of dark quarks with dark symmetry group \( SU'(2) \times SU'(N) \) and corresponding gauge bosons \( F_{\mu}^\pm \). The Yukawa couplings of the Higgs boson to the dark quark-like multiplets would obviously not align the \( SU'(2) \times SU'(N) \) with the Standard Model \( SU(3) \). In such a model it would be natural to expect the dark matter to consist of dark atoms formed from dark leptons and dark nucleons, with the atoms formed due to the \( SU'(2) \times SU'(1) \) interactions in the dark sector.

However, we will focus on the simplest model of private gauge bosons with one dark family of lepton-like particles and their private \( SU(2) \times U(1) \) gauge bosons \( U(1)_{B-L} \). This is practically equivalent to a dark multi-family model with one family being lighter than the other families, thus being the relevant family for dark matter. Because we do not need to exclude the presence of heavier lepton-like or quark-like \( SU'(2) \times SU'(1) \) doublets, we also do not have to worry about anomaly cancellation within the lightest lepton-like dark family. However, we will later find that the requirement for dark matter in the single lepton-like dark family implies \( Y_L' = 0 \), such that anomaly cancellation is actually fulfilled already with the single dark doublet \( \psi_L \).

Without dark nucleons from a dark quark sector, Coulomb-type repulsion between the dark leptons could counteract the pull of gravity. Generically this would seem to imply that the \( SU'(2) \times SU'(1) \) couplings \( q_2 \) and \( q_1 \) should only be of gravitational strength for not actually preventing dark halo formation. However, we can choose the \( SU'(2) \times SU'(1) \) charges of the dark leptons in a way which allows for the formation of \( SU'(2) \times SU'(1) \) neutral dark atoms as bound states of the dark leptons, thus alleviating the constraints on the magnitudes of \( q_1 \) and \( q_2 \).

### 3. Dark atoms

Recall that the Coulomb potential for electrons and protons in quantum optics (expressed in terms of Schrödinger picture operators)

\[
H_C = \int d^3x \int d^3x' \frac{e}{4\pi|x - x'|} \left[ \psi_+^*(x)\psi_+^*(x') \right. \\
\left. \times \psi_-(x')\psi_-(x) + \psi_+^*(x)\psi_+^*(x') \psi_-(x')\psi_-(x) \\
- \left( 2\psi_+^*(x)\psi_+^*(x') \psi_-(x')\psi_-(x) \right) \right]
\]

arises from

\[
\partial_\mu F_{\mu\nu}^\rho = e \left( \psi_+^*(x)\psi_+^*(x') - \psi_+^*(x')\psi_+^*(x) \right),
\]

(9)

in Coulomb gauge, and substitution of the solution

\[
E_{\parallel}(x) = \frac{e}{4\pi|x - x'|} \left( \psi_+^*(x')\psi_-(x) - \psi_+^*(x)\psi_-(x') \right)
\]

of Eq. (10) into the contribution from \( E_{\parallel} \) to the electromagnetic field Hamiltonian

\[
H_{EM} = \frac{1}{2} \int d^3x \left( E_{\parallel}^2(x) + E_{\perp}^2(x) + B^2(x) \right) = H_C + \frac{1}{2} \int d^3x \left( E_{\parallel}^2(x) + B^2(x) \right),
\]

(11)

\[
E_{\perp}(x) = -\partial A(x, t)/\partial t|_{t=0}, \quad B(x) = \nabla \times A(x).
\]

Substitution of the \( k \)-space expansions of the electron and proton operators and normal-ordering then yields the usual electron-electron, electron-(anti-proton), proton-proton, positron-proton etc. repulsion
between pairs of particles and attractive between particles with the charged density-density correlation operator free, e.g. if there are not sufficiently many additional dark lepton-like doublets.

In the non-relativistic limit, the equations (with $\psi^*_i \psi_i \equiv \psi^*_{\alpha\beta} \psi_{\alpha\beta}$)

$$\partial_t C^{(0)} = -q C = \frac{q}{2} Y_L^0 (\psi^*_1 \psi_1 + \psi^*_2 \psi_2) + \psi^*_{\alpha\beta} \psi_{\alpha\beta}$$

yield the dark sector Coulomb operator

$$H_{dc} = \int d^3 x \int d^3 x' \frac{q^2}{8\pi|x-x'|} \left( \sum_{a=1}^2 C_a^{(a)}(x,x') \right) \tag{12}$$

with the charged density-density correlation operator

$$C_a^{(a)}(x,x') = C_a^2(x,x') + \sum_{a=1}^2 C_a^{(a)}(x,x')$$

and anti-particles in the dark sector. The attractive channels between particles and anti-particles allow for the formation of dark $SU(2)$ mesons which will decay fast through the Higgs portal and accelerate annihilation of any remnant dark anti-leptons. We are therefore interested in attractive terms between pairs of particles for the formation of dark atoms.

The terms in Eq. (13) with coupling constants $q^2 Y_L^0 (Y_L^0 + 1)$ yield attractive interactions in the two-particle states $\psi^*_{1,L}(x)\psi_{1,R}(x')|0\rangle$ and $\psi^*_{2,L}(x)\psi_{2,R}(x')|0\rangle$ if $Y_L^0 (Y_L^0 + 1) < 0$, and attractive interactions in the two-particle states $\psi^*_{1,L}(x)\psi_{2,R}(x')|0\rangle$ and $\psi^*_{2,L}(x)\psi_{1,R}(x')|0\rangle$ if $Y_L^0 (Y_L^0 - 1) < 0$. However, the resulting bound states would have residual charges $q_{1/2}$ under $SU(2)$ and $(2Y_L^0 + 1)q_{1/2}$ under $U(1)$, respectively. This would yield repulsive Coulomb-type atom-atom interactions and again prevent dark halo collapse unless $q_1$ and $q_2$ would be tuned below gravitational strength.

However, the Coulomb term for the dark two-lepton states $\psi^*_{1,R}(x)\psi_{2,R}(x')|0\rangle$.

$$H_{1R,2R} = \int d^3 x \int d^3 x' \frac{q^2 (Y_L^0 - 1)}{16\pi|x-x'|} \times |\psi_{1,R}(x)\psi_{2,R}(x')|0\rangle \psi_{1,R}(x')|0\rangle \psi_{1,R}(x) \tag{14}$$

is attractive if $Y_L^0 < 1$. It has no $SU(2)$ charge, but a resulting atomic $U(1)$ charge $q_{1/2}$, and therefore yields uncharged dark atoms consisting of the two right-handed dark leptons. This implies overall charge neutrality of the universe also under the dark gauge groups, in the same way as anomaly cancellation in the Standard Model ensures charge neutrality in the visible sector. At temperatures above 1 TeV there are equal abundances of particle species, and the vanishing sum of visible $U(1)$ hypercharges over SM particle species and dark $U(1)$ hypercharges over dark sector species ensures overall charge neutralities in both sectors.

The potential in the dark sector therefore splits into the left-handed and right-handed parts, $H_{dc} = H_{dc} + H_{R,R}$, with fine structure constants $\alpha_L = q^2/16\pi$ and $\alpha_R = q^2/16\pi$, and the only attractive particle-particle channel for neutral atoms consisting of the two right-handed dark leptons, which are bound due to the $U(1)$ interaction. The states of the dark atoms are therefore for separation $r = x_1 - x_2$ of the two dark leptons and total atomic momentum $K$ given in the standard way by direct transcription of the corresponding results of non-relativistic QFT.

$$|\Psi_{n,f,m;K}(t)\rangle = \int d^3 x_1 \int d^3 x_2 \psi^*_{1,R}(x_1)\psi_{1,R}(x_2)|0\rangle \times |\Psi_{n,f,m;K}(x_1 - x_2)(2\pi)^{-3/2}$$
with hydrogen type wave functions $\Psi_{n,l,m}(r)$ for coupling constant $\alpha_R/4$ and reduced mass $m_{12} = m_1 m_2/M$, $M = m_1 + m_2$. The energy eigenvalues in the non-relativistic regime are

$$E(K, n) = \frac{K^2}{2M} - \frac{1}{32\pi^2} \alpha_R^2 m_{12}.$$  

Spin singlet or triplet factors were apparently omitted, since we are not interested in the fine structure of the dark atoms. The chiral projectors in the potential [14] reduce the effective dark sector fine structure constant to $\alpha_R/4$. This is explained in detail in the Appendix.

We also assume $\alpha_R < \alpha_S$ since dark $U(1)$ interactions must be weaker than electromagnetism. The separation of gas and dark matter in bullet-type clusters warrants this conclusion [28, 29]. The fact that dark matter halos are much more extended and form less concentrated substructures than baryonic matter also tells us that dark matter cannot cool down as efficiently as baryons.

4. Dark matter annihilation and thermal creation

We are interested in the non-relativistic thermal freeze-out of WIMP scale dark leptons. The annihilation of a heavy dark lepton species is then dominated by branching ratios into Standard Model particles through the Higgs portal. The $t$ and $u$ channel annihilations $\psi \bar{\psi} \to CC$ and $\psi \bar{\psi} \to XX$ are suppressed with $\alpha_R^2 p^2/m_1^2$ due to the chirality factors in the amplitudes, which come from the chirality factors $(1 \pm \gamma_5)/2$ in the vertices $q_i Y_i \bar{\psi}_i \gamma^\mu C \psi_i (1 + \gamma_5) /4$ and $q_j \bar{\psi}_j \gamma^\mu X_j \cdot \sigma (1 - \gamma_5) /4$ from Eqs. (16).

The leading order cross sections into the Standard Model states are (with $VV = W^+ W^-$ or $VV = ZZ$ and $\delta_W = 1$ for annihilation into $W^+ W^-$ or $\delta_W = 0$ otherwise)

$$\sigma_{\psi \bar{\psi} \to \gamma \gamma}(s) = \frac{1 + \delta_W}{64\pi} m_1^2 \frac{m_1^2 v}{v_s^2} \sqrt{s - 4m_1^2} \sqrt{s - 4m_1^2} \frac{(s - 2m_1^2)^2 + 8m_1^2}{(s - m_1^2)^2 + m_1^4 \Gamma_h^2}. \quad (16)$$

$$\sigma_{\psi \bar{\psi} \to f \bar{f}}(s) = \frac{g^2}{16\pi} m_1^2 m_1^2 \frac{(s - 4m_1^2)^{3/2}}{4s \sqrt{v_s^2} (s - m_1^2)^2 + m_1^4 \Gamma_h^2}, \quad (17)$$

The total annihilation cross section increases with mass $m_1$ for masses above 80 GeV, and therefore the heavier dark lepton species will determine both the mass $M$ of the dark atoms and the freeze out temperature. We will assume $m_1 \lesssim 0.01 m_2$ and therefore $M \sim m_2$.

The requirement of thermal freeze out then determines $M \sim 96$ GeV, see Fig. [1] where the thermally averaged [35] annihilation cross for a particle with mass $m_2 \sim M$ is compared to the required value from thermal dark matter creation. The logarithmically varying required value of $(\sigma v)$ for thermal creation varies from $3.38 \times 10^{-26}$ cm$^3$/s for $M = 65$ GeV to $3.45 \times 10^{-26}$ cm$^3$/s for $M = 100$ GeV.

A low mass value $m_1 \lesssim 1$ GeV implies an invisible Higgs decay width which is well below the current limits [36, 37]. $\Gamma_{h \to \psi \bar{\psi}} \lesssim 82$ keV $\sim 0.018 \Gamma_{h \to SM}$.

This dark matter model is even more predictive than the standard Higgs portal dark matter models because the coupling to the Higgs field is already determined in terms of the mass, $g = m_2/v_h \approx M/v_h$. The requirement of thermal dark matter creation therefore does not yield a parametrization $g(M)$ of the Higgs portal coupling as a function of the dark matter mass, but determines $M$. However, the corresponding nucleon recoil cross section

$$\sigma_{DN} = \frac{g^2}{48 \pi} \frac{m_1^2 M^4}{4m_1^4 v_h^2 (M + m_N)^2}, \quad (19)$$

is about $1.2 \times 10^{-44}$ cm$^2$ for $M = 96$ GeV, $m_N = 930.6$ MeV (the weighted average of the nucleon masses in
long lived xenon isotopes), and the SVZ value $g_{\nu e} = 210$ MeV \cite{38} for the effective Higgs-nucleon coupling. This is in conflict with the exclusion limits from the xenon based direct search experiments \cite{39,40}.

5. Conclusions

Alignment of gauge symmetries in the visible and dark sectors through the Higgs portal is an interesting new tool for dark matter model building. In particular, it provides a mechanism for ultraviolet complete fermionic Higgs portal models without the need of higher mass-dimension effective vertices.

Apparently, the construction presented here opens the Higgs portal into much more complicated and rich dark sectors, even with the possibility of $CP$ and time-reversal reciprocity between the visible and dark sectors, which would be broken through the different mass spectra in the two sectors. Furthermore, the construction also generalizes to alignment of gauge groups through other fields. Every field which transforms in a faithful representation of a symmetry group $G$ can align different copies $G_1, G_2, \ldots$ of the symmetry group, each with their own gauge bosons $A^a_{\mu}$ and coupling constants $g_a$, through Yukawa couplings to fields transforming under the different symmetries $G_i$. In particular, it is conceivable that the dark leptons may not couple to the Standard Model through the Higgs portal, but through a dark scalar $SU'(2)$ doublet $H'$, which couples to the dark sector gauge bosons $X_\mu$ and $C_\mu$. This would also align the electroweak-type gauge symmetries in the dark and visible sectors through the Yukawa couplings of the scalar $SU'(2)$ doublet. However, it would not be constrained by the current limits from the direct search experiments, since the Yukawa couplings of $H'$ in the visible sector must be weaker than the Higgs couplings, thus also implying a weaker effective $H'$-nucleon coupling.

Acknowledgements

This work was supported in part by the Natural Sciences and Engineering Research Council of Canada through a subatomic physics grant. I very much enjoyed the hospitality of the Kavli Institute for Cosmological Physics during my sabbatical, and I benefitted greatly from discussions with KICP members and the particle cosmology group, especially Rocky Kolb, Lian-Tao Wang, Andrew Long, Michael Fedderke, and Dan Hooper. I would also like to acknowledge Valeri Galitsky for his excellent computer support at KICP.

Appendix: The Schrödinger equation in the dark sector

To understand the impact of the chiral projectors in the dark sector Coulomb potential \cite{13}, we follow the procedure which yields the corresponding Schrödinger equation for the hydrogen atom in the baryonic sector while keeping track of the chiral projectors.

The relevant states for the dark atoms are two-particle states

$$|\Psi(t)\rangle = \sum_{\alpha\beta} \int d^3x \int d^3x' \Psi_{\alpha\beta}(x, x', t) \times \psi_{\alpha}^*(x)\psi_{\beta}^*(x')|0\rangle.$$ \hspace{1cm} (20)

The states are written in the Schrödinger picture, and the indices $\alpha, \beta$ are Dirac indices.

The relevant part of the Hamiltonian of the theory in the sector of Fock space with is spanned by the states \cite{20} (i.e. suppressing all parts of the Hamiltonian which map into different sectors of Fock space) is with $P_K = (1 + \gamma_5)/2$,

$$H = \sum_{i,\alpha,\beta} \int d^3x \bar{\psi}(x) \left( m_i \gamma_0 - i\gamma_{\beta\beta} \cdot \nabla \right) \psi(x) \times \left( P_{R\alpha} P_{R\beta} \right) \int d^3x' \int d^3x' \frac{\alpha_R}{|x - x'|} \times \psi_{\alpha}^*(x)\psi_{\beta}^*(x')\psi_{\alpha'}^*(x')\psi_{\beta'}^*(x').$$ \hspace{1cm} (21)

The corresponding Hamiltonian without the chiral projectors $P_K$ arises for the electron-proton system from the energy-momentum tensor of QED in Coulomb gauge, see e.g. Sec. 21.4 in Ref. \cite{41}.

The relativistic Schrödinger equation for the two-particle system follows from $i\hbar d|\Psi(t)\rangle/dt = H|\Psi(t)\rangle$ and after decomposition in the basis \cite{20} in the form

$$i\frac{\partial}{\partial t}\Psi_{\alpha\beta}(x, x', t) = \sum_{\rho\sigma} \tilde{H}_{\alpha\beta,\rho\sigma} \Psi_{\rho\sigma}(x, x', t),$$ \hspace{1cm} (22)

with the Hamilton operator

$$\tilde{H}_{\alpha\beta,\rho\sigma} = \left[ m_1 \gamma_0 - i(\gamma_0 \cdot \gamma)_{\alpha\beta} \cdot \frac{\partial}{\partial x} \right] \tilde{\delta}_{\rho\sigma} + \delta_{\alpha\beta} \left[ m_2 \gamma_0 \cdot \gamma_{\rho\sigma} + \left( \gamma_0 \cdot \gamma \right)_{\rho\sigma} \cdot \frac{\partial}{\partial x} \right] - \frac{\alpha_R}{|x - x'|} P_{R\alpha} P_{R\beta}.$$ \hspace{1cm} (23)

To derive the nonrelativistic limit, we use the Dirac representation of $\gamma$ matrices and write the $4 \times 4$ matrix...
\[ \Psi_{a1b}(x, x', t) \text{ in terms of } 2 \times 2 \text{ matrices in the form} \]
\[ \Psi(x, x', t) = \left( \frac{\phi(x, x', t) \phi(x, x', t)}{\xi(x, x', t) \chi(x, x', t)} \right) \times \exp\{-i(m_1 + m_2)t\}. \]

(24)

Substitution into Eq. (22) then yields in leading order of \( m^{-1} \) and \( a_R \) for the “small” components the equations
\[ \Phi(x, x', t) = -\frac{i}{2m_1} \frac{\partial}{\partial x'} \Phi(x, x', t) \cdot \sigma^T, \]
\[ \xi(x, x', t) = -\frac{i}{2m_1} \frac{\partial}{\partial x} \xi(x, x', t), \]
and \( \chi(x, x', t) = 0 \), and substitution into the equation following for \( \psi(x, x', t) \) from Eq. (22),
\[ i \frac{\partial}{\partial t} \psi(x, x', t) = \left( -\frac{1}{2m_1} \Delta - \frac{1}{2m_2} \Delta \right) \psi(x, x', t) - \frac{a_R}{4|x - x'|} \psi(x, x', t). \]

(25)

yields the standard two-particle Schrödinger equation up to an extra factor of 1/4 in the Coulomb potential,
\[ i \frac{\partial}{\partial t} \psi(x, x', t) = \left( -\frac{1}{2m_1} \Delta - \frac{1}{2m_2} \Delta \right) \psi(x, x', t) \]
\[ - \frac{a_R}{4|x - x'|} \psi(x, x', t). \]

(26)

Separation in center of mass and relative coordinates for the mapping into the effective single-particle equations with masses \( M = m_1 + m_2 \) and \( m_{12} = m_1 m_2 / M \) then proceeds as usual.

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