Flavor structures in the Dark Standard Model TeV-Paradigm

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The structure of families of dark leptons and quarks is studied as the quantum electrodynamics sector of a dark Standard Model paradigm of particle physics. We show that a minimal local U(1) symmetry with one generation of dark leptons is able to solve the problems of structure formation at small scales. Moreover, the theory of two generations provides solutions for a larger mass spectrum at the TeV-scale introducing at the same time a new energy scale of interactions in the dark sector well below the weak one.

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I. INTRODUCTION

Almost all macroscopic and microscopic natural phenomena at present times arise from a single abelian interaction between protons and electrons or even between electrons solely. One needs to look at matter at much higher energy densities to notice effects which differ from these elementary interactions. It is well known that the Standard Model (SM) is the only theory that can successfully predict almost all current experimental outcomes, once its input parameters are set.

In this work, inspired by the successes of the Standard Model, we study a low-energy secluded copy of it as a candidate for solving the dark matter (DM) problem of cosmology. This shows that the SM framework works as a universal paradigm. The effective theory that we study is self-contained and admits a straightforward UV completion. There are already a wide variety of suggested solutions to the DM problem. In particular, a lot of attention has been recently devoted to theories that alleviate the problems of ΛCDM for small-scale structure formation, i.e. the “cusp vs core”, “missing satellite” and “too big to fail” problems (for a thorough discussion see Ref. [1]). The dark Standard Model (dSM) gives clear answers to these problems.

As a low-energy limit of the dSM we consider the simple case of a broken abelian symmetry. Larger groups, like SU(2), could lead to similar results. Here, however, we neglect all other possible forces between particles in the dSM assuming that their interaction scale is much larger than the one of the abelian force. We leave the accommodating of this local abelian symmetry in a higher symmetrical theory for future works. The effective theory is thus a secluded version of quantum electrodynamics (QED) which we call dQED. Contrary to standard QED, however, we consider the possibility of a massive force mediator. This has interesting cosmological effects as explained in Ref. [2]. In addition to this massive mediator, the field content of the dQED includes \( N_\ell \) dark-lepton flavors and \( N_q \) dark-quark flavors. We show that the dQED successfully solves the small-scale problems of structure formation. We postulate that the mass of the DM candidate lies in the TeV scale. This allows for a kinetic decoupling from their relativistic scattering partners after the period of big bang nucleosynthesis (BBN), which is important for solving the small-scale problems as pointed out in Ref. [1]. Moreover, the fact that the mass of the DM candidate lies in the TeV regime is in agreement with the recent measurements of high energy electrons/positrons at the DAMPE experiment [3]. We note that there are also other dark matter models that postulate a new dark U(1) symmetry [4], [5], but they are not strictly BBN/CMB-compatible, do not study flavored DM and do not accommodate a gauge-invariant mechanism for mass generation in the UV completion [4]. While writing up this paper another model with a new dark U(1) symmetry has appeared in Ref. [7], but it involves many more new fields, since the proposed sector is not secluded from the standard model.

This paper is organized as follows. We begin with a brief motivation for the flavored dQED theory in section II. The Lagrangian and the cross sections for elastic scattering and annihilation are described in section III. The discussion of the constraints on the parameters of the theory follows in IV. Finally, we discuss in section V the cosmological observables which are relevant for solving the small-scale problems and give our results and conclusions in section VI. We will denote the dark fields with \( e \) the dark electron. To avoid confusion with the SM particles, we will rename them by adding an index SM, e.g. \( e_{SM} \) for the SM electron.

II. THE MOTIVATION

It is already known that the simultaneous solution of the enduring problems of the ΛCDM cosmology, assuming a DM candidate at the TeV scale, can be provided
by introducing at least one new force with interaction scale much below the Fermi scale [2]. For this reason we assume the existence of a hidden local U(1) symmetry (dQED) with generator $Q$ mediated by a gauge boson $X$. This is analogous to the QED paradigm. The mass of the $X$ boson is taken to lie in the MeV range [12, 8], i.e. well below the weak scale and much below the masses of the DM particles. This is necessary for a successful generation of the observed small-scale structures \([2, 9]\).

Due to lack of detection evidence, we postulate that the generation of the observed small-scale structures \([2, 9]\). This is necessary for a successful generation of the DM particles. This is analogous to the QED paradigm. The mass of the $X$ boson was in local thermal equilibrium with the SM only at early times, well before the electroweak phase transition providing masses \([2, 9]\).

For this reason we assume the existence of a hidden local U(1) symmetry \([2, 3]\). For this reason we enable chiral charges $Q_{L/R}$ for the left/right-chiral dark leptons and $Q_{L/R}$ for the quarks. The dark-Higgs charge is then fixed by

$$
\pm Q_h = Q_{L_i} - Q_{R_i} \quad \forall i \in \{1, \ldots, N_\ell\},
$$

$$
\pm Q_h = Q_{L_j} - Q_{R_j} \quad \forall j \in \{1, \ldots, N_q\}
$$

and we label its observed mass by $M_h$.

For the theory to be free of anomalies, we demand

$$
\sum_{i=1}^{N_\ell} (Q_{L_i} - Q_{R_i}) = \sum_{j=1}^{N_q} (Q_{R_j} - Q_{L_j}),
$$

$$
\sum_{i=1}^{N_\ell} (Q_{L_i}^3 - Q_{R_i}^3) = \sum_{j=1}^{N_q} (\tilde{Q}_{R_j}^3 - \tilde{Q}_{L_j}^3).
$$

The above condition together with Eqs. (1) and (2) dictates that a consistent theory should include at least one flavor of dark leptons and quarks, i.e. $N_\ell, q \geq 1$. Furthermore, $N_\ell$ and $N_q$ must be equal. This must be so, because Eq. (3) together with Eqs. (1) and (2) implies $(N_\ell \pm N_q) Q_h = 0$. For this to be satisfied we must have $N_\ell = N_q$ and $Q_{L_i} - Q_{R_i} = -(Q_{L_i} - Q_{R_i})$.

There are then two possible solutions of Eqs. (1), (2), (3) and (4) and they can be written as

$$
Q_{L_i} = \tilde{Q}_{R_i}, \quad Q_{R_i} = \tilde{Q}_{L_i}
$$

and

$$
Q_{L_i} = -\tilde{Q}_{L_i}, \quad Q_{R_i} = -\tilde{Q}_{R_i}
$$

for $i \in \{1, \ldots, N_\ell\}$. In particular this means that quarks and leptons come in pairs. From now on we will call each such pair a generation and pick the solution of Eq. (5) for the charges. Without loss of generality we assume the dark leptons to be heavier than the dark quarks.

Now we turn our attention to the technical part of the flavored dQED. Its UV-complete Lagrangian is given by

$$
\mathcal{L}_{\text{DS}} = \mathcal{L}_{\text{matter}} + \mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{Higgs}} + \mathcal{L}_{\text{Yukawa}}
$$

with

$$
\mathcal{L}_{\text{matter}} = \sum_{s \in \{L, R\}} (\bar{\ell}_s i \not{\! D} \ell_s + \bar{q}_s i \not{\! D} q_s),
$$

$$
\mathcal{L}_{\text{Yukawa}} = \bar{\ell}_L \not{\! D} \ell_R + \bar{q}_L \not{\! D} q_R + \text{H.c.},
$$

$$
\mathcal{L}_{\text{gauge}} = -\frac{1}{4} X_{\mu \nu} X^{\mu \nu}
$$

and

$$
\mathcal{L}_{\text{Higgs}} = |DH|^2 - \frac{1}{2} (|H|^2 - v_d^2)^2 + a^2 |H|^2 H_\text{SM}^1 H_\text{SM}^\dagger
$$

$\ell$ and $q$ represent the leptonic and quark generation multiplets. We have suppressed the indices $i$ of each copy. The covariant derivative is conventionally defined as $D_\mu = \partial_\mu - iX_\mu Q$ and $\not{\! D}_\mu = \not{\! D} D_\mu$ is simply

(iii) The dQED is secluded and was in local thermal equilibrium with the SM only at early times, well before big bang nucleosynthesis.

(ii) No significant new SM-philic entropy transport is present after the dark and visible sectors decouple.

(iii) For simplicity we assume that the largest contribution to the DM relic density is due to stable DM particles with mass in the TeV scale following the WIMP paradigm. No mixed DM scenarios are considered.

(iv) The field charges of the local abelian symmetry in the dQED are arbitrary, but quantum consistency is present together with a MeV interaction scale motivated by Refs. [2, 3].

III. THE FLAVORED DARK QED

The flavored dQED theory consists of $N_\ell$ dark leptons $\ell$ and $N_q$ quarks $q$, such that $N_\ell + N_q = N$. The leptons are charged with local charges $Q_\ell$, where $i \in \{1, \ldots, N_\ell\}$, and the quarks with local charges $Q_q$, where $j \in \{1, \ldots, N_q\}$. These fields serve as the source for the U(1) gauge vector boson $X$. The $X$ field is the main force mediator of the theory in the low-energy limit. As stated in assumption (iv), we require that the mass of $X$ lies in the MeV scale. The only consistent way to give $X$ such a mass $M_X$ is by a dark-Higgs mechanism. We denote the dark Higgs boson with $H$ and its local charge with $Q_h$. We assume that the dark Higgs admits a vacuum expectation value $v_d$ at times before the electroweak phase transition providing masses to the matter content of the theory. Without loss of generality we take the mass and interaction eigenstates to be identical. It is important to note that after symmetry breaking the mass of $X$ can be expressed as $M_X \sim Q_h v_d$. From this follows that $Q_h \ll Q_{\ell_i} Q_{q_j}$, since $M_X$ lies in the MeV range by assumption (iv). Therefore the dark leptons and quarks acquire only Dirac masses $m_{\ell_i}$ and $m_{q_j}$ respectively and not Majorana.
diagonal in the flavor space. The Yukawa-interaction matrix is also diagonal in the flavor space: $H_{i/q} = H/v_d \text{diag}(m^N_{i/q},...,m^N_{i/q})$. The field-strength tensor is $X_{\mu\nu} = 2\partial_{\mu}X_{\nu}$. We normalize the observed mass of the mediator and of the dark Higgs field as $M_X = \sqrt{2Q_{h}v_d}$ and $M_h = \sqrt{2\lambda v_d}$ respectively. The last term of Eq. (11) is the only dimension-4 contact between the dQED and the SM which is allowed by the symmetries with $\alpha$ being real.

IV. CONSTRAINTING THE PARAMETER SET

In this section we present possible constraints on the parameters of the dQED theory. The first constraint comes from requiring partial-wave unitarity of the scattering matrix (see Ref. [10]). This leads to an upper bound $m^2_h < O(10^3)$ TeV for the mass of the leptons acting as the DM candidate. Secondly we look at the constraints on the couplings: the couplings must be small enough, to allow a perturbation expansion, since we are interested in the QED paradigm at low energies. The effectiveness of the interactions between the Standard Model and the dark sector depends strongly on the values of the dark Higgs mass and the bridge coupling $\alpha$. Bounds on the coupling $\alpha$ can be found in [11]. We remark that, after assuming $v_d \gtrsim M_h \sim$ TeV, $\alpha \lesssim 10^{-2}$ is phenomenologically expected.

The DM candidates are thermally produced WIMPs. We define $g_i$ as $g_i \equiv \frac{1}{2} (Q_{Li} + Q_{Ri})$. From now on we make the approximation $g_i \approx Q_{Li} \approx Q_{Ri}$, which is valid to lowest order in $Q_h \ll g_i$. Assuming a single generation of dark leptons and quarks, the corresponding WIMP condition on the coupling constant $\alpha$ is

$$4\pi\alpha_{ij} > \left(\frac{m^2_{\ell /q}}{M^4_{\text{Pl}}} \right)^{1/2},$$

where $x_i \equiv m_i/T_i$ is the normalized inverse temperature at the $i$th-DM freeze out and $4\pi\alpha_{ij} \equiv g_i g_j$. This condition ensures that the dark leptons were in thermal equilibrium with the dark radiation before they started annihilating. The generalization to more generations is straightforward.

The deviation of the effective number of degrees of freedom of SM neutrinos $\Delta N_{\text{eff}}$ can be modified by the presence of lighter dark leptons and quarks. $\Delta N_{\text{eff}}$ parametrizes the relativistic energy density at different times of the cosmological evolution. Its deviation due to the presence of the above discussed dark sector is

$$\Delta N_{\text{eff}}|_{\text{BBN}} = N_\nu \frac{\rho_\eta}{\rho_\nu},$$

where $\rho_\eta$ is the energy density of the relativistic dark species. The values at the SM-neutrino decoupling temperature $T_{\nu D} = 2.3$ MeV [12] are taken as initial conditions. We define

$$\varepsilon|_{\nu} := \left(\frac{T_{DS}}{T_{\nu}}\right)^3 = g^*_{DS}(T_D) g^*_{SM}(T_{\nu D}) g^*_{DS}(T_{\nu}),$$

where $g^*_{DS}$ are the entropy degrees of freedom of the dark sector and $g^*_{SM}$ are the ones of the Standard Model. Assumption [10] has been used to derive this expression. For the expected values of $\alpha$ the dark sector is no longer in local thermal equilibrium with the SM plasma around the temperature of the electroweak phase transition, $T_{\text{EW}}$. A value of $\varepsilon|_{\text{BBN}} = 0.288$ is found for $N = 1$ taking $m_\ell \sim$ TeV and $m_q \lesssim T_{\text{BBN}}$, or equivalently for $N = 2$ with $m^2_1 \sim m^2_2 \sim$ TeV, $m^2_3 \sim T_{\text{EW}}$ and $m^2_4 \lesssim T_{\text{BBN}}$. The results are 1σ-compatible with BBN-data [13] and CMB-data [14] and correspond to $\Delta N_{\text{eff}}|_{\text{BBN}} \approx 0.38$. It is important to note that there is a relation between $\Delta N_{\text{eff}}$ and the mass of the lighter dark quarks $m_q$. We consider for simplicity a single quark species with $m_q < T_{\text{BBN}}$. These dark quarks are the last scattering partners of the DM candidates and determine the kinetic decoupling temperature $T_{kd}$. Since a late kinetic decoupling has positive consequences for the solution of the small-scale problems, the lighter dark quarks should still be relativistic or semi-relativistic at that time. This means roughly $m_q \lesssim 16\text{ TeV}$. Furthermore, the contribution of these particles to the energy density of the universe after they freeze-out is estimated with $\Omega_\eta h^2 \approx m_q/255\text{ eV}$, since we assume that they are stable and that the dark sector is secluded. As stated in assumption [10], a mixed DM ensemble for the relic density is unfavored. For this reason, we take conservatively $\Omega_\eta h^2$ to be less than 0.0245 [14], which corresponds to a minimum amount of mixed DM relic density. This gives an upper bound $m_q \lesssim 0.65\text{ eV}$. In addition, assuming that the recombination takes place instantaneously at 0.3 eV, the energy contribution of these light fermions at this time yields

$$\frac{\Delta N_{\text{eff}}|_{\text{CMB}}}{\varepsilon|_{\nu}^{1/3}} = 240 \frac{\pi^4}{x_{\text{CMB}}^3} \sqrt{\frac{z^2}{\exp(z) + 1}},$$

where $x_{\text{CMB}} = m_q (4/11 \varepsilon|_{\text{CMB}})^{-1/3}/T_{\text{rec}}$. The impact on $\Delta N_{\text{eff}}$ is maximized for $m_q \sim T_{\text{rec}}$. For example, if $m_q = 0.6\text{ eV}$ then $N_{\text{eff}} = 3.14$, which is in agreement with the value obtained from the Planck measurements [14]. Such values of $m_q$ may also explain the recent tension about the decrease of the deviation of effective SM-neutrino number from BBN-based to CMB-based measurements. Such a difference, $\Delta N_{\text{eff}}|_{\text{CMB}} - \Delta N_{\text{eff}}|_{\text{BBN}} < 0$, arises naturally from this theory.

V. COSMOLOGICAL OBSERVABLES

In this section we discuss the most important cosmological observables for a DM candidate: (A) the relic abundance, (B) the SIDM virtue, and (C) the characteristic damping scales. In particular we indicate the optimal values that these cosmological observables should
attain and their dependence on the parameters of the theory. An explicit benchmark point for the parameters is then presented in Section VI.

A. The relic abundance of flavors

A condition on the couplings of the theory can be found by requiring that the theory yields the correct DM relic density. By assumption (iii) the DM population consists mostly of dark leptons. Each dark lepton with mass $m_\ell^i$ annihilates rapidly for $T < m_\ell^i$ through the dominant $s$-wave channels of the theory if the spectrum hierarchy allows it. The approximated cross section is

$$\frac{m_\ell^i}{m_\ell^i} (\langle \sigma \tau \rangle_{\text{rel}})_i \approx \Theta(m_\ell^i - M_X) \alpha_i^2 + \sum_{i \neq j} \Theta(m_\ell^i - m_\ell^j) \alpha_{ij}^2.$$ \hspace{1cm} (16)

$\langle ... \rangle$ denotes the thermal average and the last sum over $j$ runs over all dark fermions. The first term in Eq. (16) is due to annihilations into gauge bosons and the second one to annihilations into lighter fermions. We assumed that $M_h > m_\ell^i$, because we are not interested in annihilations into dark-Higgs bosons, which are $p$-wave suppressed. The above annihilation cross section is a first order approximation, because it admits corrections of order $(Q_h m_\ell^i/v_d)^2$.

We do not consider them in our analysis, since under the natural assumption $m_\ell^i \lesssim v_d$ such contributions are subdominant. It is straightforward to show that the thermal evolution of the theory is not modified up to $Q_h^2$ corrections to the observables. Therefore, without loss of generality we can assume $M_h \gtrsim m_\ell^i$ and neglect annihilations into SM particles after the electroweak phase transition. In order to obtain an expression for the relic density, we start from the Boltzmann equation for the DM distribution function and solve it numerically following [15]. The resulting relic abundance per particle/antiparticle is

$$\Omega_{\chi} h^2 \approx 0.06 \left( \frac{\alpha_i(N)}{0.13} \right)^{-2} \left( \frac{m_\ell^i}{10 \text{ TeV}} \right)^2.$$ \hspace{1cm} (17)

using the abbreviation

$$\alpha_i(N) = \sqrt{\Theta(m_\ell^i - M_X) \alpha_{ii}^2 + \sum_{i \neq j} \Theta(m_\ell^i - m_\ell^j) \alpha_{ij}^2}.$$ \hspace{1cm} (18)

This result includes the Sommerfeld effect [16], which gives corrections of order $\mathcal{O}(1)$ towards smaller couplings. Note that the self-annihilation cross sections are well below the experimental sensitivity [17].

B. Self-interactions of flavors

The dark sector includes self-interaction processes. Appropriate values of the corresponding cross sections have been shown to solve the “cusp vs core” and “too big to fail” problems. See for example Refs. [2, 13, 19] for an explanation. The required cross-section values are $(\langle \sigma \tau / m_\ell^i \rangle_{\text{therm}} \sim 1 \text{ cm}^2 \text{g}^{-1}$ at the scale of dwarf galaxies and $(\langle \sigma \tau / m_\ell^i \rangle_{\text{therm}} \sim 0.1 \text{ cm}^2 \text{g}^{-1}$ at the scale of clusters [2]. In these expressions the cross sections are averaged over a Maxwell-Boltzmann thermal distribution where $v_{\text{therm}}$ is the most probable velocity. The dQED naturally includes self-interactions between the dark matter candidates. It can thus be regarded as a true self-interacting dark matter (SIDM) theory. The averaged cross-section for self-interaction per unit mass, $(\langle \sigma \tau / m_\ell^i \rangle_{\text{therm}}$, is strongly velocity-dependent. Numerical solutions of the momentum transfer cross section in SIDM theories with non-relativistic fermions can be found in Refs. [3, 20]. We use the classical results given in Ref. [21]. Combining the effects of self-interactions with a late kinetic decoupling of dark matter has positive consequences for the solution of the “cusp vs core” and “too big to fail” problems, as shown in Refs. [22, 23].

C. Damping masses

The chemical decoupling of DM candidates from the dark radiation plasma, which was discussed in \textsection A, does not determine the moment of last contact between the two. The elastic scatterings between the heavy leptons $i$ and the lighter quarks $j$ at late times constitute the longest-running channel for efficient momentum exchange in the dark sector. The $i$-species remain in kinetic equilibrium with the light quarks longer than with the gauge bosons if one assumes $M_X \gg m_q^i$. At temperatures $M_X \gg T \gg m_q^i$ and at lowest order in perturbation theory, the averaged momentum-transfer elastic cross section per lepton flavor is

$$\langle v_{\text{rel}} \sigma_{T} \rangle_{ij} \approx \frac{40 \xi(3) G_{ij}^2 T^2}{\pi \zeta(3)} \left( \frac{T_{\text{DS}}}{T} \right)^2.$$ \hspace{1cm} (19)

$T$ is the average photon temperature and $T_{\text{DS}}$ the dark sector temperature. $\sigma_T$ is defined as $\sigma_T \equiv \int \text{d} \Omega (1 - \cos \theta) |d \sigma_{\text{el}} / d \Omega |$. and $G_{ij} \equiv \sqrt{2} g_i g_j / 4 M_X^2$ is the Fermi constant associated to the corresponding interaction. In the above expression we ignored the quark mass.

Before kinetic decoupling, the effective elastic interactions damp perturbations in the linear power spectrum, which would otherwise grow and form the first DM structures, i.e. protohalos [24]. After $T_{\text{kd}}$ the interactions are too weak to keep sustaining local thermal equilibrium and the remaining elastic scatterings can be described as sources of entropy in an imperfect DM fluid [25]. The corresponding damping masses can be estimated by $M_d = (4\pi/3) \rho_m(T_{kd}) / \mathcal{H}^3(T_{kd})$, since all the DM candidates within the Hubble radius at kinetic decoupling were in thermal contact with the dark radiation plasma. In general, collisionless damping or free streaming should be taken into account as well. This is not the
case here, because for values of \( m \) in the TeV range free streaming is negligible. As stated in Refs. [1, 23, 26], the “missing satellite” problem can be solved by damping masses of the order of dwarf galaxies, i.e. approximately \( 10^{8-9} M_\odot \). This is equivalent to \( T_{kd} \sim 1 \text{ keV} \) [2], which satisfies the Lyman-\( \alpha \) bound, as stated in Refs. [1, 23, 27, 28] and more recently in Ref. [29]. These values are not excluded by any current collider and DM direct-search constraints and can be accommodated very well in this model. We point out that CDM theories usually obtain damping masses in the range of the earth mass [30], which lies well below the resolution capabilities of current numerical simulations [31]. In order to satisfy Lyman-\( \alpha \) measurements [28], we consider a temperature of kinetic decoupling \( T_{kd} \) after BBN and before the sub-keV epoch, as in Ref. [15]. More recent constraints are presented in Ref. [24], but they might be overly restrictive. An analytic derivation of the expression for the temperature of kinetic decoupling \( T_{kd} \) can be found in Ref. [22]. Applying this prescription to the case \( N_\ell = 1 \) or equivalently for any \( N \) as long as the relic abundance is fixed through the freeze out of the \( i \)-th dark lepton, we obtain

\[
T_{kd} \approx 147 \text{eV} \epsilon_{\ell k}^{-1/2} \left( \frac{m_\ell}{10 \text{ TeV}} \right)^{-1/4} \left( \frac{\Omega_\ell h^2}{0.12} \right)^{1/4} \left( \frac{Q_\ell}{2 \times 10^{-8}} \right) \left( \frac{v_\ell}{40 \text{ TeV}} \right) \left( \frac{T_{\nu}}{T_{\nu}^0} \right)^{3/2}.
\]

Note that this result does not depend directly on \( \alpha(N) \), which is fixed by [17]. Changes in \( \Omega_\ell \) do not play a big role either. The strongest dependence of \( T_{kd} \) is the one with respect to the dark-Higgs charge and its vacuum expectation value.

**VI. RESULTS AND CONCLUSION**

We now present a choice of parameters of the dQED theory solving the small-scale problems of \( \Lambda \text{CDM} \). As a benchmark point we can choose for \( N_\ell = 1 \) thermally produced WIMPs with a mass \( m_\ell = 10 \text{ TeV} \) and a mediator mass of \( M_X = 1.1 \text{ MeV} \). This corresponds to a kinetic decoupling temperature of \( T_{kd} \approx 0.456 \text{ keV} \), yielding the correct damping masses. The related sections for self interactions are then in the desired range and similar to those of the ETHOS 4 model [23]. The case of \( N_\ell = 2 \), however, appears more flexible (for example following the spectrum hierarchy proposed in section IV), since two different generation couplings are present and even masses around 2 TeV seem to solve the small- and large-scale problems simultaneously.

This flavored dQED paradigm enables also rare annihilations into SM particles. For example, if we assume \( M_h \approx 2 m_\ell \), the \( \ell \)-leptons can annihilate through a resonant channel into an SM-electron/positron pair. This can be interpreted as a cosmic ray excess (CRE) at the TeV scale. In particular the velocity-averaged cross sections after the electroweak phase transition can be as large as \( \sim 10^{-20} \text{cm}^2 \text{s}^{-1} \). This is in line with the DAMPE-CRE measurements [3]. However, we note that the dark leptons can also decay into hadrons. Therefore, this construction alone cannot explain the observed cosmic ray peak without some modifications. One possible solution is to introduce a specific leptophilic interaction [32]. Such phenomenological extensions of the dQED theory are beyond the scope of this paper. In particular, we note that for \( M_h \approx 2 m_\ell \) rare annihilations into SM-particles consistent with the recent measurements of excesses in the cosmic electron flux [3] are possible. This is in agreement with the LEP constraints discussed in Ref. [34]. However, a precise description of the needed additional leptophilic interactions is beyond the scope of this paper.

In this work we studied the flavor structure of an abelian quantum gauge symmetry which is UV-complete and called it dQED. This dQED could be interpreted as the low-energy part of a dark standard model. In particular, we showed that this theory satisfies all cosmological and particle physics constraints and solves the small-scale problems of structure formation in a rather simple way compared a neurinophilic theory [15]. Even a purely phenomenological model, which explains the long standing anomaly of the neutron lifetime in the DM context [35], is partially based on the dQED prescription. In order to construct the theory, we introduced a SM-singlet family of flavored DM consisting of light and heavy fields and we found that one should introduce a new energy scale of interactions below the weak one in order to obtain positive cosmological signatures. However, it should be clear that even larger groups as for example SU(2) could straightforwardly lead to similar results and, therefore, we only considered the dQED theory as a starting point. The fact that this very simple model can comfortably accommodate all the important constraints is very promising.

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