The dark photon is a new gauge boson whose existence has been conjectured. It is dark because it arises from a symmetry of a hypothetical dark sector comprising particles completely neutral under the Standard Model interactions. Dark though it is, this new gauge boson can be detected because of its kinetic mixing with the ordinary, visible photon. We review its physics from the theoretical and the experimental point of view. We discuss the difference between the massive and the massless case. We explain how the dark photon enters laboratory, astrophysical and cosmological observations as well as dark matter physics. We survey the current and future experimental limits on the parameters of the massless and massive dark photons together with the related bounds on milli-charged fermions.

CONTENTS

I. Introduction
A. Massless and massive dark photons
1. Kinetic mixing: Electric or hyper-charge? 4
2. Embedding in a nonAbelian group 5
B. UV models
1. Massive dark photon: Origin and size of the mixing parameter 5
2. Massless dark photon: Higher-order operators 6
C. Dark matter and the dark photon
1. Massless dark photon and galaxy dynamics 7
2. Massless dark photon and dark-matter relic density 7
3. Massive dark photon and light dark matter 8
4. Massive dark photon as dark matter 9

II. Phenomenology of the massless dark photon
A. Limits on the dark dipole scale $d_{M}/\Lambda^{2}$ 10
1. Astrophysics and cosmology 10
2. Precision, laboratory and collider physics 12
3. Can the massless dark photon be seen at all? 15
B. Limits on milli-charged particles 15
C. A minimal model of the dark sector 17

III. Phenomenology of the massive dark photon
1. Production, decays and detection 21
2. Visible and invisible massive dark photon 22
A. Limits on the parameters $\varepsilon$ and $m_{A'}$ 23
1. Constraints for $m_{A'} > 1$ MeV with $A'$ decays to visible final states 23
2. Constraints for $m_{A'} > 1$ MeV with $A'$ decays to invisible final states 26
3. Constraints for $m_{A'} < 1$ MeV 27
B. Constraints on $y$ and $m_{\chi}$ 29

IV. Concluding remarks 31
Acknowledgments 32
A. Dark sector portal 32
B. Boltzmann equation and relic density 33
C. Thermal field theory 34
References 35
I. INTRODUCTION

New particles beyond the Standard Model (SM) have always been thought to be charged under at least some of the same gauge interactions of ordinary SM particles. This assumption has driven the theoretical speculations as well as the experimental searches of the last 50 years but it has also been increasingly challenged by the negative results of all these searches—and the mounting frustration for the failure to discover any of these hypothetical new particles.

As the hope of a breakthrough along these lines is waning, interest in a dark sector—dark because not charged under the SM gauge groups—is growing. Maybe no new particles have been seen simply because they do not interact through the SM gauge interactions.

The dark sector may contain few or many states, and these can be fermions or scalars or both, depending on the model. Dark matter proper is found among these states, whose relic density can be computed and constrained by observational data. In addition, these dark states can interact; their interactions can be Yukawa-like or mediated by dark gauge bosons or both depending on the model.

If the dark and the visible sectors were to interact only gravitationally—which they cannot avoid—there would be little hope of observing in the laboratory particles belonging to the dark sector. A similar problem exists for the dark sector in general, and the dark photon in particular: The main motivation in introducing new-physics scenarios is to use them as a foil for the SM in mapping possible experimental discrepancies. In the absence of clearly identified new states, the many parameters, for instance, of the minimal supersymmetric SM or even of the effective field theory approach to physics beyond the SM, are working against their usefulness. Instead, each dark sector can be reduced to few parameters—to wit, just two in the case of the dark photon—in terms of which the possible discrepancies with respect to the SM are more effectively mapped in the experimental searches and the potential discovery more discernible.

As explained in detail in section I.A, there are actually two kinds of dark photons: The massless and the massive—whose theoretical frameworks as well as experimental signatures are quite distinct. They give rise to dark sectors with different features; their characteristic physics and experimental searches are best reviewed separately. The massive dark photon has been receiving so far most of the attention because it couples directly to the SM currents and is more readily accessible in the experimental searches. The massless dark photon arises from a sound theoretical framework and, as we shall argue, provides, with respect to the massive case, a comparably rich, if perhaps more challenging, experimental target.

We look into the ultraviolet (UV) completion of these models in section I.B, and, more in general, in (Georgi et al., 1983; Okun, 1982)

The concept of a portal—which at first blush might seem rather harmless—actually represents a radical departure from what is the main conceptual outcome of our study of particle physics, namely, the gauge principle and idea that all interactions must be described by a gauge theory. The portal, and the new interactions that it brings into the picture, adds a significant exception to this principle. Among the possible portals, the vector case deviates the least from the gauge principle as it only introduces a mixing for the gauge bosons while the interaction to matter remains of the gauge type (albeit with an un-quantized charge). Instead, the other kinds of portal imply a manifest new violation to the gauge principle, the other’s being the notable case of the Yukawa and self-interactions of the Higgs boson—which are themselves, exactly because of their not being gauge interactions, the least understood part of the SM.

There is an additional and important reason to study the dark sector in general, and the dark photon in particular: The main motivation in introducing new-physics scenarios is to use them as a foil for the SM in mapping possible experimental discrepancies. In the absence of clearly identified new states, the many parameters, for instance, of the minimal supersymmetric SM or even of the effective field theory approach to physics beyond the SM, are working against their usefulness. Instead, each dark sector can be reduced to few parameters—to wit, just two in the case of the dark photon—in terms of which the possible discrepancies with respect to the SM are more effectively mapped in the experimental searches and the potential discovery more discernible.

The portal consists of relevant operators that can take various forms depending on the spin of the mediator: Vector (spin 1), Neutrino (spin 1/2), Higgs (scalar) and Axion (pseudo-scalar) portals are the best motivated and most studied (see appendix A).

Among these possible portals, the vector portal is the one where the interaction takes place because of the kinetic mixing between one dark and one visible Abelian gauge boson (non Abelian gauge bosons do not mix). The visible photon is taken to be the boson of the $U(1)$ gauge group of electromagnetism—or, above the electroweak symmetry-breaking scale, of the hyper-charge—while the dark photon\(^1\) comes to be identified as the boson of an extra $U(1)$ symmetry (an idea first considered in the context of supersymmetric theories in (Fayet, 1980, 1981) and, more in general, in (Georgi et al., 1983; Okun, 1982)).

\(^1\) The names para- (Holdom, 1986b), hidden-sector, secluded photon and U-boson (Fayet, 1990) have also being used to indicate this particle.
section III. At the best of our knowledge, these two sections provide the reader with a comprehensive review of the physics of the dark photon.

We collect in three appendices a number of definitions and equations, which the reader may find useful to better follow the discussion in the main text.

In the past few years a number of reports on the dark sector (and the massive dark photon within it) have been published (Alekhin et al., 2016; Alexander et al., 2016; Beacham et al., 2020; Cucuárello, 2016; Deliyergiyev, 2016; Essig et al., 2013; Hewett et al., 2012; Raggi and Kozhuharov, 2015). The interested reader can therein find different points of view to complement the present review as well as additional details on the other portals.

A previous discussion of the astrophysical, cosmological and other constraints for the massless dark photon can be found in (Dobrescu, 2005).

A. Massless and massive dark photons

It is useful to identify from the very beginning two kinds of dark photons:

- the massless kind, which, as we are about to show, does not couple directly to any of the SM currents and interacts instead with ordinary matter only through operators of dimension higher than four;
- the massive kind, which couples to ordinary matter through a current (with arbitrary charge), that is, a renormalizable operator of dimension four. The massless limit of this case does not correspond to the massless case above.

Because of their different coupling to SM particles, the two kinds are best discussed separately.

Let us first consider the massless case. As first discussed in (Holdom, 1986b) in this case the classical Lagrangian can be diagonalized. What happens at the quantum level and how the mixing manifests itself has been analyzed in detail in (del Aguila et al., 1995) for the unbroken gauge theory as well as the spontaneously broken case (see, also, the appendix of (Feldman et al., 2007) which we mostly follow).

The most general kinetic part of the Lagrangian of two $U(1)_a$ and $U(1)_b$ gauge bosons is

$$\mathcal{L}_0 = -\frac{1}{4} F_{\mu\nu}^a F_a^{\mu\nu} - \frac{1}{4} F_{\mu\nu}^b F_b^{\mu\nu} - \frac{\varepsilon}{2} F_{\mu\nu}^a F_{b}^{\mu\nu}. \quad (1)$$

The gauge boson $A_\mu^a$ is taken to couple to the current $J_\mu$ of ordinary SM matter, the other, $A_\mu^b$, to the current $J'_\mu$, which is made of dark-sector matter:

$$\mathcal{L} = e J_\mu A_\mu^a + e' J'_\mu A_\mu^b, \quad (2)$$

with $e$ and $e'$ the respective coupling constants.

![FIG. 1 Scheme of the coupling of the ordinary ($A_\mu$) and dark ($A'_\mu$) photon to the SM and dark-sector (DS) particles for the two choices of the angle $\theta$ discussed in the main text. $e$ and $e'$ are the couplings of the ordinary and dark photons to their respective sectors.](image)

The kinetic terms in Eq. (1) can be diagonalized by rotating the gauge fields as

$$\begin{pmatrix} A_\mu^a \\ A_\mu^b \end{pmatrix} = \begin{pmatrix} 1 & 0 \\ -\frac{1}{\sqrt{1-\varepsilon^2}} & 1 \end{pmatrix} \begin{pmatrix} \cos \theta & \sin \theta \\ -\sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} A_\mu^a \\ A_\mu^b \end{pmatrix}, \quad (3)$$

where now we can identify $A_\mu^a$ with the ordinary photon and $A_\mu^b$ with the dark photon. The additional orthogonal rotation in Eq. (3) is always possible and introduces an angle $\theta$ which is arbitrary as long as the gauge bosons are massless.

After the rotation in Eq. (3), the interaction Lagrangian in Eq. (2) becomes

$$\mathcal{L}' = \left[ e' \cos \theta J'_\mu + e \left( \sin \theta - \frac{\varepsilon \cos \theta}{\sqrt{1-\varepsilon^2}} \right) J_\mu \right] A_\mu^a$$

$$+ \left[ -e' \sin \theta J'_\mu + e \left( \cos \theta + \frac{\varepsilon \sin \theta}{\sqrt{1-\varepsilon^2}} \right) J_\mu \right] A_\mu^b. \quad (4)$$

By choosing $\sin \theta = 0$ ($\cos \theta = 1$) (see right-side of Fig. 1) we can have the ordinary photon $A_\mu^a$ coupled only to the ordinary current $J_\mu$ while the dark photon couples to both the ordinary and the dark current $J'_\mu$, the former with strength $\varepsilon \varepsilon / \sqrt{1-\varepsilon^2}$ proportional to the mixing parameter $\varepsilon$. The Lagrangian is therefore:

$$\mathcal{L}' = \left[ e' \cos \theta J'_\mu - \frac{\varepsilon \varepsilon}{\sqrt{1-\varepsilon^2}} J_\mu \right] A_\mu^a + e J_\mu A_\mu^b. \quad (5)$$

Vice versa, with the choice $\sin \theta = \varepsilon$ and $\cos \theta = \sqrt{1-\varepsilon^2}$ (see left-side of Fig. 1), we have the opposite situation with the dark photon only coupled to the dark current and the ordinary photon to both currents, with strength $\varepsilon \varepsilon / \sqrt{1-\varepsilon^2}$ to the dark one. This latter coupling between the dark-sector matter to the ordinary photon is called a milli-charge. Its value is experimentally known to be small (Davidson et al., 2000). The dark photon sees ordinary matter only through the effect of operators...
like the magnetic moment or the charge form factors (of dimension higher than four). This is the choice defining the massless dark photon proper:

\[ \mathcal{L}' = e' J'_\mu A^\mu + \left[ -\frac{e' \varepsilon}{\sqrt{1 - \varepsilon^2}} J'_\mu + \frac{e}{\sqrt{1 - \varepsilon^2}} J_\mu \right] A^\mu \]  

(6)

If the gauge symmetry is spontaneously broken, the diagonalization of the mass terms locks the angle \( \theta \) to the value required by the rotation of the gauge fields to the mass eigenstates and we cannot have that one of the two currents only couples to one of the two gauge bosons.

This is also the case when the \( U(1) \) gauge bosons acquire a mass by means of the Stueckelberg Lagrangian (see Ruegg and Ruiz-Altaba, 2004) for a review and the relevant references)

\[ \mathcal{L}_{Stu} = -\frac{1}{2} M^2 \delta a \delta A^\mu_a - \frac{1}{2} M^2 \delta b \delta A^\mu_b - M a b \delta A^\mu_{a b} \]  

(7)

In this case, as in the spontaneously broken case, the angle \( \theta \) is fixed and equal to

\[ \sin \theta \equiv \frac{\delta \sqrt{1 - \varepsilon^2}}{\sqrt{1 - 2 \delta \varepsilon + \delta^2}} \quad \cos \theta \equiv \frac{1 - \delta \varepsilon}{\sqrt{1 - 2 \delta \varepsilon + \delta^2}} \]  

(8)

where \( \delta = M_b / M_a \), and we have no longer the freedom of rotating the fields as in Eq. (3). The Lagrangian in Eq. (4) is now

\[ \mathcal{L}'' = \frac{1}{\sqrt{1 - 2 \delta \varepsilon + \delta^2}} \left[ \frac{e' (1 - \delta \varepsilon)}{\sqrt{1 - \varepsilon^2}} J'_\mu + \frac{e (\delta - \varepsilon)}{\sqrt{1 - \varepsilon^2}} J_\mu \right] A^\mu + \frac{1}{\sqrt{1 - 2 \delta \varepsilon + \delta^2}} \left[ e J_\mu - e' J'_\mu \right] A^\mu. \]  

(9)

The case of spontaneously broken symmetry can be distinguished from the Stueckelberg mass terms because the former will give rise to processes in which the dark photon is produced together with the dark Higgs boson, the vacuum expectation value of which hides the symmetry.

Whereas the Lagrangian in Eq. (9) is the most general, the simplest and most frequently discussed case consists in giving mass directly to only one of the \( U(1) \) gauge bosons so that, for instance, \( M_b = 0 \) in Eq. (7), the mass states are already diagonal. Even in this simple case, the mass term removes the freedom of choosing the angle \( \theta \) in Eq. (3). With this choice, \( \delta = 0 \) in Eq. (9), the ordinary photon couples only to ordinary matter and the massive dark photon is characterized by a direct coupling to the electromagnetic current of the SM particles (in addition to that to dark-sector matter) and described by the Lagrangian

\[ \mathcal{L} \supset -\frac{e \varepsilon}{\sqrt{1 - \varepsilon^2}} J_\mu A^\mu \sim -e J_\mu A^\mu, \]  

(10)

as in Eq. (5) above. This is the choice defining the massive dark photon. The coupling of the massive dark photon to SM particles is not quantized—taking the arbitrary value \( e \varepsilon \). Because of this direct current-like coupling to ordinary matter, it is the spontaneously broken or Stueckelberg massive dark photon that is mostly discussed in the literature and considered in the experimental proposals.

Notice that the massive dark photon has the same couplings as the massless dark photon after choosing \( \sin \theta = 0 \) (right-side of Fig. 1); this case therefore represents the limit of vanishing mass of the massive dark photon. On the contrary, the massless dark photon proper—corresponding to the choice \( \tan \theta = [\varepsilon / \sqrt{1 - \varepsilon^2}] \)—is not related to any limiting case of the massive dark photon.

There are no electromagnetic milli-charged particles in the massive case; they are present only if both \( U(1) \) gauge groups are spontaneously broken (or equivalently \( M_b \neq 0 \) in the Stueckelberg Lagrangian in Eq. (7))—which is not the case of our world where the photon is massless.

1. Kinetic mixing: Electric or hyper-charge?

There seems to be the choice in the kinetic mixing in Eq. (1) between the \( U(1)_{e.m.} \) group of electric charge and the \( U(1)_Y \) group of the hyper-charge, with mixing parameter \( \varepsilon \) defined as in Eq.(1). Concerning the massless dark photon, these two choices give rise to the same physics, since the dark photon remains decoupled from the SM fields at the tree-level. The only difference is that the photon and Z-boson are now both coupled to the dark-sector current, with \( e_e \varepsilon / \sqrt{1 - \varepsilon^2} \cos \theta_W \) and \( e_e \varepsilon / \sqrt{1 - \varepsilon^2} \sin \theta_W \) strength, respectively.

Let us now consider the massive dark-photon coupling to hyper-charge. In this case it is convenient to parametrize the coupling of the dark photon to the hyper-charge as

\[ \mathcal{L} = -\frac{\varepsilon}{2 \cos \theta_W} \tilde{E}^{\mu \nu} B_{\nu \nu}. \]  

(11)

The usual diagonalization of the gauge bosons \( W^3_\mu \) and \( B_\mu \) now includes also the dark photon \( A'_\mu \) (in the non-diagonal basis) so that the physical gauge bosons \( Z_\mu \) and \( A'_\mu \) also contain a dark-photon component \( A'_\mu \) in the mass eigenstate basis. In particular, at the \( O(\varepsilon) \) in the expansion, we have

\[ \begin{pmatrix} W^3_\mu \\ B'_\mu \\ A'_\mu \end{pmatrix} \cong \begin{pmatrix} c_W & s_W & -s_W \varepsilon \\ -s_W & c_W & -c_W \varepsilon \\ t_W & e & 1 \end{pmatrix} \begin{pmatrix} Z_\mu \\ A'_\mu \\ A'_\mu \end{pmatrix}, \]  

(12)

where \( c_W, s_W, \) and \( t_W \) are the usual cosine, sine, and tangent of the Weinberg angle \( \theta_W \), respectively. New couplings of the massive dark photon to the SM fermions appear for the photon and the Z gauge boson up to \( O(\varepsilon^2) \):

\[ \mathcal{L} \supset -e J^\mu A'_\mu + e' e t_W J^\mu Z_\mu + e' J^\mu A'_\mu, \]  

(13)
where \( J_\mu \) is the EM current, while \( J'_\mu \) and \( \epsilon' \) are the matter current and coupling of the dark-photon in the dark sector, respectively. After integrating out the \( Z \) boson, we see that the coupling of the massive dark photon to the SM fermions is recovered as \(-\epsilon \epsilon'\).

Which coupling is used depends then only on the energy of the processes considered, with the direct coupling to the photon for all processes below the electroweak scale breaking, and the hyper-charge above it. Since all limits are to be considered approximately within the order of magnitude, the presence of the factor \( c_W \) in the definition in Eq. (11) does not matter. The Lagrangian in Eq. (13) shows that, if the mixing is between the dark photon and the hyper-charge, the \( Z \) gauge boson acquires a milli-charged coupling strength \( \epsilon' t_W \epsilon' \) to the dark sector current.

For completeness, let us also recall two other possibilities that have been discussed in the literature:

- There is no kinetic mixing as in Eq. (1) but the mass term between the dark photon and the \( Z \)-boson is taken non-diagonal and therefore giving a mixing between these two states (Appelquist et al., 2003; Babu et al., 1998; Davoudiasl et al., 2012; Galison and Manohar, 1984; He et al., 1991). The dark photon is named the dark \( Z \) and there are characteristic experimental signatures in parity violating processes and the coupling to neutrinos;

- The \( B - L \) global symmetry (or other conserved flavor symmetries) are gauged and taken to be the \( U(1) \) group of the dark photon, which mixes with the hyper-charge (Bauer et al., 2018; Fayet, 2017; Hooft, 2014). There is direct coupling to the SM fermions in this case and the dark photon is no longer dark.

Although their implementation is not discussed in this review, other interesting generalizations—as, for instance, the dark photon to be considered a Kaluza-Klein state in a model with large extra-dimensions (Rizzo, 2018) or the interplay between the neutrino see-saw mechanism and the dark photon (Bertuzzo et al., 2019)—should be borne in mind.

### 2. Embedding in a nonAbelian group

In the massless case, the ordinary photon still couples to the dark sector with a milli-charge \( \epsilon \epsilon' \). As reviewed in the next section, there are very stringent limits on the size of such a milli-charge, at least for reasonably light dark states. To avoid the necessity of assuming a very small milli-charge, one can assume that the dark \( U(1) \) group is a symmetry left over after the spontaneous breaking of a larger nonAbelian group.

The simplest realization of this symmetry breaking is provided by the group \( SU(2) \) spontaneously broken to \( U(1) \) by the vacuum expectation value of the neutral component of a scalar field in the adjoint representation.

In this scenario, the mixing term in Eq. (1) cannot be written because the larger group has traceless generators. The absence of mixing is in this case protected against radiative corrections and the dark and the ordinary photons see only their respective sectors (at least through renormalizable operators).

This scenario is also suggested by the extra Landau pole that otherwise would be present—assuming that the Landau pole of the ordinary \( U(1) \) is removed by the embedding of the SM in a scenario of grand unified theory.

If we assume that the dark photon arises from a non-Abelian group, there is no milli-charged coupling of the dark sector to ordinary photons. On the other hand, all states in the dark sector must come as multiplets of the nonAbelian group and the possible experimental signatures of this additional structure can be searched for.

### B. UV models

Because the massive dark photon couples directly to the SM electromagnetic current, its phenomenology is rather independent of the details of the underlaying UV completion. The two parameters \( \epsilon \) and \( m_{A'} \) suffice to fully describe the experimental searches.

The case of the massless dark photon is more complicated because the coupling to the SM particles only takes place through higher order operators whose structure heavily depends on the underlaying UV model. Even though it is possible to frame the experimental search in terms of the effective scale of these operators (as we do in section II), the limits thus found begs to be framed in terms of the UV model parameters, namely the masses and the coupling of the dark sector states, in addition to the dark photon itself. For this reason, it is useful in this case to introduce a minimal UV model (as we do in section II.C) to provide the relationships among the parameters of the model and thus possible to relate different limits that are instead independent or not present under the portal interaction.

#### 1. Massive dark photon: Origin and size of the mixing parameter

The size of the mixing parameter \( \epsilon \) is arbitrary. It is this feature that makes the charge not quantized. At the same time, it cannot be \( O(1) \) because, if so, the massive dark photon would have already been discovered.

A natural suppression of \( \epsilon \) is achieved if the mixing only comes as a correction at one- or two-loop level in some UV completion. This is achieved in a natural manner if the tree-level mixing is set to zero. One looks for the renormalization of the model and introduces the nec-
cessary counter-terms, of which the mixing in Eq. (1) is one. If there are states in the UV completion carrying both ordinary and dark charges, the loop of these states generates the mixing but it comes suppressed by the loop factor (neglecting logarithmic terms) and therefore of order, say, $1/(16\pi^2)$ times the square of the coupling constant and therefore approximately $O(10^{-3})$, for a perturbative value of such a coupling. One can further suppress such a term by assuming that the states carrying both charges come in doublets of opposite dark charges. In this case, the first contribution is at the two-loop level, and approximately of order $O(10^{-5})$. If the mixing originates in the exchange of heavy messenger fields (Grzadkowski et al., 2010) or in a multi-loop contribution (Gherghetta et al., 2019; Koren and McGehee, 2020), its value can be smaller.

Even smaller values of the parameter $\varepsilon$ are expected if the origin of the mixing is non-perturbative; for example, values between $O(10^{-12})$ and $O(10^{-6})$ have been discussed—mostly within the broad heading of string compactification (Abel and Schofield, 2004; Abel et al., 2008; Dienes et al., 1997; Goodsell, 2009; Goodsell et al., 2009; Heckman and Vafa, 2011), or in scenarios of SUSY breaking (Arkani-Hamed and Weiner, 2008) and hidden valley (Chan et al., 2012). These arguments are often cited to motivate experimental searches in the region of small mixing parameter $\varepsilon$ in the case of the massive dark photon—regardless of the large uncertainties in the predictions of the corresponding theoretical approaches.

2. Massless dark photon: Higher-order operators

The massless dark photon does not interact directly with the currents of the SM fermions. The higher-order operators through which the interaction with ordinary matter $\psi'$ takes place start with the dimension-five operators in the Lagrangian

$$\mathcal{L} = \frac{\varepsilon_D}{\Lambda^2} \bar{\psi} i \gamma_\mu \sigma_{\mu\nu} \left( D^{ij}_M + i \gamma_5 D^{ij}_E \right) \psi^j F^{\mu\nu},$$

(14)

where $F^{\mu\nu}_D$ is the field strength associated to the dark photon field $A'_\mu$, and $\sigma_{\mu\nu} = i/2 [\gamma_\mu, \gamma_\nu]$. The operator proportional to the coefficient $D_M$ is the magnetic dipole moment and that proportional to the coefficient $D_E$ is the electric dipole moment. The indices $i$ and $j$ in the fermion fields keep track of the flavor and thus allow for flavor off-diagonal transitions.

The dimension-five operators in Eq. (14) are best seen as operators of dimension six with the gauge group $SU(2)_L$ taken as the unbroken symmetry of the Lagrangian and the SM fermion grouped, like in the SM, into doublets $\psi_L$ and singlets $\psi_R$. In this case, the operators contain the Higgs boson field and can be written as

$$\mathcal{L} = \frac{\varepsilon_D}{2\Lambda^2} \bar{\psi}_L i \gamma_\mu \sigma_{\mu\nu} \left( D^{ij}_M + i \gamma_5 D^{ij}_E \right) H \psi^j_R F^{\mu\nu} + \text{H.c.}$$

(15)

The effective scale is accordingly modulated by the vacuum expectation value (VEV) $v_h$ of the Higgs boson. This VEV keeps track of the chirality breaking, with the whole operator vanishing as $v_h$ goes to zero.

In this review we shall only retain the magnetic dipole $D_M$ term and set to zero the electric dipole term proportional to $D_E$. The inclusion of the latter would require the further assumption of CP-odd physics which is, we believe, premature at the moment.

Next, we have the dimension-six operators

$$\mathcal{L}' = \frac{\varepsilon_D}{2\Lambda^2} \bar{\psi} i \gamma_\mu \gamma_5 \left( R^{ij}_M + i \gamma_5 R^{ij}_a \right) D_\mu \psi^j F^{\mu\nu},$$

(16)

where the form factor $R_M$ is related to the charge radius of the fermion; the term $R_a$ is sometimes referred to as the anapole.

The operator in Eq. (16) contributes, via the equations of motion, to four-fermion operators—which are accounted for in the effective field theory of the dimension-six operators (Grzadkowski et al., 2010) but are not relevant for the massless dark photon interaction to ordinary matter—and to the form factors of the interaction if the particles are off-shell. The latter provide a next-to-leading interaction between the massless dark photon and ordinary matter that has yet to be studied (and is not discussed in this review).

Higher-order operators give vanishingly small contributions and can be neglected.

The scale $\Lambda$ depends on the parameters of the underlying UV model. Typically, it is the mass of a heavy state, or the ratio of masses of states of the dark sector, multiplied by the couplings of these states to the SM particles. In particular, the dipole operators in Eq. (15), as they require a chirality flip, can turn out to be enhanced, or suppressed, according to the underlying model chirality mixing.

The fact that the interaction between the massless dark photon and the SM states only takes place through higher-order operators provide an appealing explanation...
for its weakness. The structure of these operators leads directly to the possible underlaying UV models—a minimal example of which is discussed in section II.C.

C. Dark matter and the dark photon

Dark matter is part of the dark sector. The interplay between the dark photon and dark matter opens new windows on its physics and gives further constraints. Whereas in most scenarios dark matter is one of the fermion (or scalar) states in this sector, there also exists the possibility that dark matter could be a very light vector boson like the massive dark photon itself.

1. Massless dark photon and galaxy dynamics

Models of self-interacting dark matter charged under Abelian or non-Abelian gauge groups and interacting through the exchange of massless as well as massive particles have a long history.\(^2\)

The most obvious obstacle to having dark matter in the dark sector interacting via a long-range force as the one carried by the massless dark photon comes from the essentially collisionless dynamics of galaxies and the ellipticity of their dark-matter halo.

The most severe observational limits come from the present dark matter density distribution in collapsed dark matter structures, rather than effects in the early Universe or the early stages of structure formation (Ackerman et al., 2009; Cyr-Racine and Sigurdson, 2013; Feng et al., 2009).

Bounds have been derived from the dynamics in merging clusters, such as the Bullet Cluster (Clowe et al., 2006), the tidal disruption of dwarf satellites along their orbits in the host halo, and kinetic energy exchanges among dark matter particles in virialized halos. The latter turns out to be the most constraining bound, noticing that self-interactions tend to isotropize dark matter velocity distributions, while there are galaxies whose gravitational potentials show a triaxial structure with significant velocity anisotropy; limits have been computed, with subsequent refinements, via estimating an isotropization timescale (through hard scattering and cumulative effects of many interactions, also taking into account Debye screening) and comparison to the estimated age of the object (Feng et al., 2009), or following more closely the evolution of the velocity anisotropy due to the energy transfer (Agrawal et al., 2017). The ellipticity profile inferred for the galaxy NGC720, according to (Agrawal et al., 2017) sets a limit of about

\[
m_X \left( \frac{0.01}{\alpha_D} \right)^{2/3} \gtrsim 300 \text{ GeV},
\]

where \(m_X\) stands for the dark matter mass and the \(\alpha_D\) scaling quoted is approximate and comes from the leading \(m_X\) over \(\alpha_D\) scaling in the expression for the isotropization timescale.

The limit in Eq. (17) is subject to a number of uncertainties and assumptions; it is less stringent than earlier results, such as the original bound from soft scattering quoted in (Ackerman et al., 2009),

\[
\frac{G_N m_X^4 N}{8 \alpha_D^2} \gtrsim 50 \log \frac{G_N m_X^2 N}{2 \alpha_D^2},
\]

where \(N\) is the number of dark-matter particles and \(G_N\) is Newton’s constant, as well about a factor of 3.5 weaker than (Feng et al., 2009) (see, also, (Feng et al., 2010; Liu et al., 2012)). On the other hand, results on galaxies from \(N\)-body simulations in self-interacting dark matter cosmologies (Peter et al., 2013), which take into account predicted ellipticities and dark matter densities in the central regions, seem to go in the direction of milder constraints, about at the same level or slightly weaker than the value quoted in Eq. (17)—again subject to uncertainties, such as the role played by the central baryonic component of NGC720.

2. Massless dark photon and dark-matter relic density

All the stable fields within the dark sector provide a multicomponent candidate for dark matter whose relic density depends on the value of their couplings to the dark photons and SM fermions (into which they may annihilate, depending on the UV model) and masses.

Not all of the dark fermions contribute to the relic density. If these fermions are relatively light, their dominant annihilation is into dark photons (see Fig. 2)

\[
\chi \chi \rightarrow A' A'
\]

with a rate given by

\[
\langle \sigma_{\chi \chi \rightarrow A' A' v} \rangle = \frac{2 \pi \alpha_D^2}{m_X^2}.
\]

For a strength \(\alpha_D \simeq 0.01\), all fermions with masses up to around 1 TeV have a large cross section and their relic

---

\(^2\) The literature on the subject is already very extensive, see, for example, (van den Aarssen et al., 2012; Ackerman et al., 2009; Agrawal et al., 2017; Arkani-Hamed et al., 2009; Baldi, 2013; Boddy et al., 2014; Buckley and Fox, 2010; Buen-Abad et al., 2015; Carbon et al., 1992; Chu and Dasgupta, 2014; Cline et al., 2014, 2012; Cyr-Racine and Sigurdson, 2013; Feng et al., 2009, 2008; Foot, 2004; Gabrielli and Raidal, 2014; Goldberg and Hall, 1986; Gradwohl and Frieman, 1992; Holdom, 1986a; Hooper et al., 2012; Kaplan et al., 2010; Tulin et al., 2013).

Interacting dark matter can form bound states. The phenomenology of such atomic dark matter (Kaplan et al., 2010) has been discussed in the literature, see (Cyr-Racine and Sigurdson, 2013) and references therein.
density (see Eq. (B15) in the appendix B)

\[ \Omega_\chi h^2 \approx \frac{2.5 \times 10^{-10} \text{ GeV}^{-2}}{\langle \sigma v \rangle} \]  \hspace{1cm} (21)

is only a percent of the critical one; it is roughly $10^{-4}$ the critical one for dark fermions in the 1 GeV range, even less for lighter states. These dark fermions are not part of dark matter; they have (mostly) converted into dark photons by the time the universe reaches our age and can only be produced in high-energy events. This is fortunate because, as we have seen, they are ruled out as possible dark matter candidates by the limit on galaxy dynamics.

Heavier dark fermions can be dark matter. The dominant annihilation via these is not into dark photons but into SM fermions via the exchange of some messenger field $S$—the details depending on the underlying UV model—and is proportional to the corresponding coupling which we denote $\alpha_L$ anticipating the discussion in section II.C—with a thermally averaged cross section approximately given by

\[ \langle \sigma \chi \chi \rightarrow ff \rangle \approx \frac{2\pi \alpha_T^2}{m_S^2} \]  \hspace{1cm} (22)

instead of Eq. (20). The critical relic density can be reproduced if, assuming thermal production,

\[ 2\pi \alpha_L^2 \left( \frac{10 \text{ TeV}}{m_S} \right)^2 \approx 0.1. \]  \hspace{1cm} (23)

These dark matter fermions belonging to the dark sector are in principle detectable through the long range exchange of the massless dark photon and its coupling to the magnetic (or electric) dipole moment of SM matter which is induced at the one loop level in the UV model of the dark sector. The somewhat complementary problem of dark matter having dipole moment and interacting with nuclei through the exchange of a photon has been discussed in (Banks et al., 2010; Barger et al., 2011; Chu et al., 2020; Del Nobile et al., 2012; Fornengo et al., 2011; Pospelov and ter Veldhuis, 2000; Sigurdson et al., 2004) This dipole interaction is now included within the basis of the operators in the effective field theory of dark matter detection (Brod et al., 2018; Fitzpatrick et al., 2013; Liem et al., 2016).

3. Massive dark photon and light dark matter

When dark matter is lighter than the dark photon, and $m_{A'} > 2m_\chi$, the annihilation channel (see Fig. 2)

\[ \chi \chi \rightarrow A' \rightarrow ff \]  \hspace{1cm} (24)

is open and we are in a scenario with light dark matter (LDM) (Boehm et al., 2004; Essig et al., 2012b; Knapen et al., 2017). The cross section is

\[ \sigma_{\chi \chi \rightarrow ff} = \frac{4\pi \varepsilon^2 \alpha_p m_\chi^2}{3} \left( 1 + \frac{2m_\chi^2}{s} \right) \left( 1 + \frac{2m_\chi^2}{s} \right) \frac{1}{(s - m_{A'}^2)^2 + m_{A'}^2 \Gamma_{A'}^2} \sqrt{1 - \frac{4m_\chi^2}{m_{A'}^2}} \]  \hspace{1cm} (25)

and dark matter was originally discussed in (Boehm and Fayet, 2004) and, more recently, in (Hambye et al., 2019). It has been suggested (Izaguirre et al., 2015) that the best variable to plot most effectively the constraints in the case of LDM is by means of the yield variable

\[ y \equiv \varepsilon^2 \alpha_p \left( \frac{m_\chi}{m_{A'}} \right)^4 \]  \hspace{1cm} (28)

because, from Eq. (26)

\[ \langle \sigma \chi \chi \rightarrow ff \rangle \approx \frac{16\pi \alpha y}{m_\chi} \]  \hspace{1cm} (29)

and therefore the relic density is brought into the plot. Moreover, the scaling of these limits is made less depen-
Massive dark photon as dark matter

A very light massive dark photon could be a dark-matter candidate if produced non-thermally in the early Universe as a condensate, the same way as the axion is produced by the misalignment mechanism (Abbott and Sikivie, 1983; Dine and Fischler, 1983; Preskill et al., 1983). In this mechanism, the value of the field is frozen by the fast expanding Universe to whatever value it has at the initial moment. The rate of expansion is much larger than the mass and the field has no time to relax to the minimum of the potential. The unavoidable (and troublesome) Lorentz-invariance violation is estimated to be small and undetectable.

In this scenario for the dark photon, as discussed in (Arias et al., 2012; Nelson and Scholtz, 2011), the mass arises via the Stueckelberg mechanism and there must be a non-minimal coupling to gravity. Once the Hubble constant value drops below the mass of the dark photon, its field starts to oscillate and these oscillations behave like non-relativistic matter, that is, like cold dark matter.

There exist two constraints on the parameters of this dark photon scenario. First of all, the initial value must be fine-tuned to reproduce the critical density. Second, the decay into photons and SM leptons must not affect the cosmic microwave background. This latter requirement means that the mixing parameter $\epsilon$ must not be too large (roughly, less than $10^{-9}$) and the mass $m_{A'}$ must be less than 1 MeV.

Production by fluctuations during inflation provides another possibility of having a massive dark photon as dark matter (Graham et al., 2016; Nakai et al., 2020).

The dark-photon dark matter is non-relativistic and interacts with ordinary matter mostly through the photo-electric process in which a photon (with energy $\omega$) is captured by an atom, with atomic number $Z$, with a cross section given, for ordinary photons, by

$$\sigma_{\text{p.e.}} = 4\alpha^4\sqrt{2}Z^8 \frac{8\pi r_e^2}{3} \left(\frac{m_e}{\omega}\right)^{7/2},$$

where $\omega$ is the photon energy and $r_e$ the classical radius of the electron. The cross section for the dark photons is that of ordinary photons rescaled by the mixing parameter $\epsilon$:

$$\sigma_{A'} = \epsilon^2 \sigma_{\text{p.e.}}.$$  

This scenario is made accessible to the experiments by considering the rate of absorption of the dark photon by the detector (Bloch et al., 2017; Pospelov et al., 2008):

$$\Gamma_{A'} = \frac{\rho_{A'}}{m_{A'}} \sigma_{A'} v_{A'},$$

where the density $\rho_{A'}$ is estimated from the relic density (or the flux from the Sun).

II. PHENOMENOLOGY OF THE MASSLESS DARK PHOTON

The phenomenology of the massless dark photon depends on the effect of the higher-order operator in
Eq. (15) which mediates its interaction with the SM particles. This operator enters the measured observables with an effective scale $\Lambda$ and the absolute value

$$d_{ij}^M \equiv |D_{ij}^M|$$

(36)

of the magnetic dipole coefficient (neglecting the CP-odd $D_E$) which can eventually be related to the parameters of the underlying UV model like masses and coupling constants. The experimental searches can thus be framed in terms of the scale $\Lambda$, the dipole coefficient $d_{ij}^M$ and and the dark charge coupling $e_D$, which we rewrite as $\alpha_D = e_D^2/4\pi$. We do not assume this scale and coefficient to be universal. Depending on the particular experimental set-up, the constraints are further sensitive to which particular lepton or quark is actually taking part in the interaction. The index, or indices, $i$ and $j$ keep track on the flavor dependence.

We discuss in section II.B the other side of the massless dark photon, namely the search for dark particles coupled to the ordinary photon by a milli-charge.

A. Limits on the dark dipole scale $d_M/\Lambda^2$

We collect in this section the known constraints on the size of the operator in Eq. (15).

We show in Fig. 4 and 5 the more stringent limits. Though these limits are on the combinations $d_M/\Lambda^2$, with a factor depending on $\alpha_D$, we find it convenient to plot them as $d_M$ as a function of $\Lambda$ so as to easily see what values of the dipole coefficient are allowed given a value for the scale $\Lambda$ (and two representative value of $\alpha_D$).

1. Astrophysics and cosmology

Astrophysics and cosmology provide very stringent limits on the interaction of the dark photon with SM matter as given by the operator in Eq. (15). It is understood that all the limits are mostly on the order of magnitude because of intrinsic uncertainties in the astrophysics of stellar medium, supernova dynamics and cosmological processes.

Astrophysical constraints for models with a massless dark photon can be derived from those obtained for axion-like particles because the dipole operator in Eq. (15) gives, in the non-relativistic limit, a derivative (and spin-dependent) coupling of the dark photon with momentum $k$ and polarization $\epsilon$ to ordinary fermions $\psi$ given by

$$\mathcal{M}_{A'_\mu} \approx \bar{\psi} (k \times \epsilon) \cdot \sigma \psi,$$

(37)

which, after averaging over the polarizations, gives the same contribution as that for a pseudo-scalar particle (Hoffmann, 1987; Raffelt, 1996) like the axion, namely

$$\mathcal{M}_a \approx \bar{\psi} k \cdot \sigma \psi.$$

(38)

Only a factor of two must be included for the independent polarizations of the dark photon.

Because the massless dark photon does not mix with the ordinary photon, we can compute the limits in a kinetic theory in which the amplitude for the relevant process is computed in the vacuum and the effect of the medium—be it the stellar interior or the supernova nucleon gas—is included in the abundances of the SM states at the given temperature.

- **Stars.** The luminosity of stars is related to their energy balance. This balance is a sensitive probe of the stellar dynamics and the particle-physics processes on which is based. Three processes are important for energy loss in stars: Compton scattering, pair creation and Bremsstrahlung. Of these three, it is the latter that provides the most stringent limit. The non-observation of anomalous energy transport, in various different types of stars, places strong constraints on the dipole coupling between SM states and the dark photon (Carlson, 1987; Dobrescu, 2005).

The quantity we need is the energy loss due to the emission of the extra particle. The energy loss per unit volume $Q$ is given in the appendix C in terms of the squared amplitude of the process of emitting, in our case, an axion.

For the Bremsstrahlung emission of axions, by electrons in the field of $n$ nuclei with charge $Z_j$, the squared amplitude is (Nakagawa et al., 1987; Raffelt, 1990)

![FIG. 3 Bremsstrahlung of dark photons from electrons in a star and from nucleons in a supernova.](image-url)
where $p_1$ and $p_2$ and $k$ are the momenta of the initial electrons and $q = p_2 - p_1$. $\omega$ and $k$ the energy and momentum of the axion and $\kappa_F = (4\sqrt{p_F E_F} / \pi)^{1/2}$ where $p_F$ and $E_F$ are the Fermi momentum and energy of the electrons in the plasma. The coefficient $\alpha_{ae}'$ is the coupling constant of the axion to the electrons.

In a degenerate medium (like the one for red giants and white dwarves) we have that the energy-loss rate per unit mass $Q / \rho$ is given by (Raffelt, 1996)

$$Q / \rho = \frac{\pi^2 \alpha^2 \alpha_{ae}'}{15 m_e} \sum_{\text{spin}} Z_j^2 n_j T^4 \sum_{\text{j}} Z_j^2 n_j F(\kappa_F) \simeq \alpha_{ae}' 1.08 \times 10^{27} \left( \frac{T}{10^9 K} \right)^4 \frac{\alpha^2}{\Lambda} F(\kappa_F),$$

the latter equation is written in units of erg $g^{-1}$s$^{-1}$, and the factor $F$ is approximately given in the relativistic limit as

$$F(\kappa_F) \simeq \frac{2 + \kappa_F^2}{2 \ln \frac{2 + \kappa_F^2}{\kappa_F^2} - 1}.$$  

The most stringent limit for electrons comes from cooling in white dwarves (Miller Bertolami et al., 2014) and giant red stars (Viaux et al., 2013) by axion Bremsstrahlung in a degenerate medium. A combined fit of the data (Giammotti et al., 2016) finds (at 2$\sigma$) that the coupling must be

$$\alpha_{ae}' \leq 3.0 \times 10^{-27}.$$  

The bound in Eq. (42) is translated into a bound for the dark photon by identifying the combination of parameters in the operator in Eq. (15) that controls the same process. This correspondence yields the equation

$$\alpha_{ae}' = 2 \frac{1}{4\pi} \left(2e_D d_M^\alpha v_h m_e / \Lambda^2\right)^2,$$

where the factor of 2 in front takes into account the two polarizations of the dark photon (with respect to the axion), $v_h = 174$ GeV and $m_e$ is the electron mass.

To satisfy the limit in Eq. (42), the dark photon parameters in Eq. (43) must satisfy

$$\frac{\Lambda^2}{\sqrt{\alpha_{ae}' d_M^\alpha}} \gtrsim 4.5 \times 10^6 \text{ TeV}^2,$$

after having included the numerical values of $m_e$ and $v_h$.

The limit in Eq. (44) updates the one found in (Dobrescu, 2005).

- Supernovae. An additional limit is found from the neutrino signal of supernova 1987A, for which the length of the burst constrains anomalous energy losses in the explosion.

As before, a bound can be derived from that for the coupling between axions and nucleons. The corresponding averaged square amplitude is given in (Brinkmann and Turner, 1988; Raffelt and Seckel, 1995) as

$$Q / \rho \simeq \alpha_{aN}' 1.74 \times 10^{13} \frac{\rho}{10^{15}} \left( \frac{T}{\text{MeV}} \right)^6,$$

in units of erg $g^{-1}$s$^{-1}$, which should not exceed the neutrino luminosity. This limit yields, taking the most conservative estimate in (Carenza et al., 2019; Keil et al., 1997),

$$\alpha_{aN}' \leq 1.3 \times 10^{-18}.$$  

The combination that controls energy transfer to dark photons in this process from ordinary matter (the quarks...
in the nucleons) is
\[ \alpha'_{sN} = 2 \frac{1}{4\pi} \left( 2 e_D d_M^2 \frac{\nu_h m_N}{\Lambda^2} \right)^2, \tag{49} \]
where \( m_N \) is the nucleon mass. By taking the limit in Eq. (48), we have
\[ \frac{\Lambda^2}{\sqrt{\alpha_D d_M^2}} \gtrsim 4.3 \times 10^5 \text{ TeV}^2, \tag{50} \]
which applies to the light \( u \) and \( d \) quarks—if we neglect small corrections due to the form factors in going from the nucleons to the quarks. The limit in Eq. (50) updates the one found in (Dobrescu, 2005).

A caveat in the limit in Eq. (48) is due to the fact that if the coupling is too strong the emitted axions are re-absorbed by the expanding supernova and there is no energy loss; this happens for
\[ \alpha_{sN}' \geq 0.7 \times 10^{-14}, \tag{51} \]
which yields
\[ \frac{\Lambda^2}{\sqrt{\alpha_D d_M^2}} \lesssim 5.9 \times 10^3 \text{ TeV}^2, \tag{52} \]
There are however limits from laboratory physics, discussed in the next section, that almost close this window.

- **Big bang nucleosynthesis.** A cosmological bound for the dark photon operator in Eq. (15) comes from the determination of the effective number of relativistic species in addition to those of the SM partaking in the thermal bath—the same way the number of neutrinos is constrained. This number is constrained by data on big bang nucleosynthesis (BBN) to be (Fields et al., 2020):
\[ N_{\text{eff}} = 2.878 \pm 0.278. \tag{53} \]

We follow (Dobrescu, 2005) in deriving the corresponding limits.

The two degrees of freedom of the dark photon exceeds this limit at the big bang temperature \( T_{\text{BBN}} \) and must have decoupled before at temperature \( T_d \) which is taken to be just above the QCD phase transition: \( T_d = 150 \) MeV. The request of decoupling before the BBN epoch can be translated in having the Hubble constant (see appendix B)
\[ H(T_d) = \frac{T_d^2}{M_{\text{Pl}}} \left( \frac{\pi^2}{90} g_*(T_d) \right)^{1/2}, \tag{54} \]
be larger than the rate of interactions between SM states and the dark photon
\[ \Gamma_{A'} = n_{A'} \langle \sigma v \rangle, \tag{55} \]
where \( \langle \sigma v \rangle \) is the thermally averaged cross section for the interaction of the dark photon with the SM particles present at the temperature \( T_d, v = 1 \), and the number density of dark photon is given (see appendix B) by \( (k_B = 1) \)
\[ n_{A'} = \frac{2 \zeta(3)}{\pi^2} T^3. \tag{56} \]

The cross section for SM fermions to Compton and annihilate into dark photon is approximately given by
\[ \langle \sigma v \rangle \simeq \frac{\alpha_D d_M^2 v_h^2}{\Lambda^4}. \tag{57} \]

We thus find the condition
\[ \frac{2 \zeta(3)}{\pi^2} T_d^4 \langle \sigma v \rangle < \frac{T_d^2}{M_{\text{Pl}}} \left( \frac{2\pi^2}{45} g_*(T_d) \right)^{1/2}, \tag{58} \]
where the effective number of degrees of freedom \( g_*(T_d) \) is bound from the limit on \( N_{\text{eff}} \). This relationship is obtained from
\[ \left( \frac{T_{\text{BBN}}}{T_d} \right)^4 = \left( \frac{g_*(T_{\text{BBN}})}{g_*(T_d)} \right)^{4/3} < \frac{7}{4} \Delta N_{\text{eff}}, \tag{59} \]
which, knowing that \( g_*(T_{\text{BBN}}) = 43/4 \), gives
\[ g_*(T_d) > (43/7)^{4/3} \Delta N_{\text{eff}}^{-3/4}, \tag{60} \]
where \( \Delta N_{\text{eff}} \equiv N_{\text{eff}} - 3 \simeq 0.34 \) by taking \( 2\sigma \) of the result in Eq. (53).

The limit applies to the interaction of leptons (electron and muon):
\[ \frac{\Lambda^2}{\sqrt{\alpha_D d_M^2}} \gtrsim 6.6 \times 10^3 \text{ TeV}^2, \tag{61} \]
and quarks \((u, s, d)\):
\[ \frac{\Lambda^2}{\sqrt{\alpha_D d_M^2}} \gtrsim 4.3 \times 10^3 \text{ TeV}^2, \tag{62} \]
which partake into the Compton and annihilation processes. The difference between Eq. (61) and Eq. (62) is due to the number of colors.

2. Precision, laboratory and collider physics

Laboratory physics can set new constrains on the dipole operator in Eq. (15). They are less stringent than those from astrophysics and cosmology because the higher-order dipole operator always yields a small number of events; these small numbers are amplified in the stars by the enormous density of particles in the medium but not in the laboratory experiments where the density is smaller.
• Precision physics. The operator in Eq. (15) gives rise to a macroscopic spin-dependent (non-relativistic) potential (Dobrescu and Mocioiu, 2006):

\[
V(r) = -\frac{\alpha_D v^2 d_M^a d_M^b}{4\Lambda^4 r^4} \left[ \sigma_a \cdot \sigma_b - 3 (\sigma_a \cdot \hat{r}) (\sigma_b \cdot \hat{r}) \right],
\]

(63)
where \( \mathbf{r} = \mathbf{r}_a - \mathbf{r}_b \) is the vector distance and \( r = |\mathbf{r}| \) and \( \hat{r} \) the corresponding unit vector. The potential in Eq. (63) is between two fermions \( f_a \) and \( f_b \), with spin \( \sigma_a \) and \( \sigma_b \), and magnetic dipole moments \( d_{M[a]}^f \), as defined in Eq. (15)—whose interaction can affect atomic energy levels as well as macroscopic forces.

The potential in Eq. (63) can be used to explore atomic physics as well as macroscopic fifth-force like interactions.

Many atomic energy levels are known with high precision. Unfortunately, the theoretical computation is lagging behind many of the experiments, mainly because of uncertainties in higher-order corrections like those due to the size of the nuclei. For this reason many results are given as energy differences where corrections proportional to \( 1/r^3 \) are factorized out. This procedure makes often impossible to use these results to test the potential in Eq. (63).

The best limit is obtained in the fine-structure spectroscopy of Helium. The extra interaction between the two electrons has been discussed in (Ficek et al., 2017) whose limits, obtained by the constraints from the \( 2^3P_2-2^3P_1 \) transitions in He, can be expressed as

\[
\frac{\Lambda^2}{\sqrt{\alpha_d d_M^p}} \gtrsim 872 \text{ GeV}^2. \tag{64}
\]

Bounds on long-range forces depending on spin set limits on the scale of the operator in Eq. (14) based on the potential in Eq. (63) as discussed in (Dobrescu and Mocioiu, 2006). The strongest bounds come from limits on macroscopic forces between electrons (Ni et al., 1999)

\[
\frac{\Lambda^2}{\sqrt{\alpha_d d_M^p}} \gtrsim 1.61 \text{ TeV}^2, \tag{65}
\]

and electrons and nucleons (Wineland et al., 1991)

\[
\frac{\Lambda^2}{\sqrt{\alpha_d d_M^p}} \gtrsim 1.94 \text{ TeV}^2. \tag{66}
\]

The limits among nucleons and electrons and protons are weaker.

Whereas the strong limits on the anomalous magnetic momentum of the electron and the muon are traditionally used to set limits on new physics, they cannot be used directly in our case because they only apply to operators coupling to the visible photon. The operator in Eq. (14) enters in the computation of the magnet moments but only at higher order with two insertions in the loop computation. The limits are accordingly weak. The contribution of the dark photon to the anomalous magnetic moment is given by

\[
a_{\ell \ell}^A = -\frac{3}{2} \frac{\alpha_D}{\pi} \left( \frac{m_\ell v_h d_M^p}{\Lambda^2} \right)^2 \left[ \frac{5}{4} + \log \left( \frac{\Lambda^2}{m_\ell^2} \right) \right] \tag{67}
\]

in the \( \overline{MS} \) scheme; contrary to the SM case, the result depends on the subtraction of a divergence.

We discuss below in section II.C how in the UV model, where there are states coupled to both the dark and the visible photon, the anomalous magnetic moment can be brought to bear directly on the limits.

The quantity \( \Delta a_e \), the difference between the experimental value of the electron anomalous magnetic moment (Hanneke et al., 2008) and its SM prediction is very small. The uncertainty on this difference (at 1\( \sigma \)) is given by (Gindice et al., 2012)

\[
\delta \Delta a_e < 8.1 \times 10^{-13}. \tag{68}
\]

By requiring that the contribution of the dark photon does not exceed this value, and therefore does not contribute to the electron magnetic moment, we obtain

\[
\frac{\Lambda^2}{\sqrt{\alpha_d d_M^p}} \gtrsim 0.075 \text{ TeV}^2, \tag{69}
\]

by taking the renormalization scale \( \mu = m_\ell \).

The analogous quantity \( \Delta a_\mu \), the difference between the experimental value of the muon anomalous magnetic moment (Bennett et al., 2006) and its SM prediction (Blum et al., 2018), is less than

\[
\Delta a_\mu < 2.74 \times 10^{-9}, \tag{70}
\]

at 2\( \sigma \) level, from which we derive

\[
\frac{\Lambda^2}{\sqrt{\alpha_d d_M^p}} \gtrsim 0.5 \text{ TeV}^2, \tag{71}
\]

for \( \mu = m_\mu \). Notice that the current 3.2\( \sigma \) discrepancy in \( \Delta a_\mu \) could be explained by requiring

\[
\frac{\Lambda^2}{\sqrt{\alpha_d d_M^p}} \simeq 0.27 \text{ TeV}^2. \tag{72}
\]

Flavor changing processes can provide constraints on possible dipole operator contributions to off-diagonal interactions.

In the lepton sector the process \( \mu \rightarrow eX^0 \), with \( X^0 \) a massless neutral boson, is bounded to (Bayes et al., 2015)

\[
\text{BR}(\mu \rightarrow eX^0) < 5.8 \times 10^{-5}, \tag{73}
\]

which gives

\[
\frac{\Lambda^2}{\sqrt{\alpha_d d_M^p}} \gtrsim 5.1 \times 10^5 \text{ TeV}^2. \tag{74}
\]

Similar limits in the hadron sector on, for example, the decays \( K \rightarrow \pi X^0 \) or \( B \rightarrow KX^0 \), cannot be used because they are forbidden when \( X^0 \) is a spin one boson like the dark photon. The decay \( B \rightarrow K^+X^0 \) is not forbidden but gives a very weak bound. Instead, the current limit on the rare decay \( K^+ \rightarrow \pi^+\nu\bar{\nu} \) given by (at the 90\% CL, see, for example, (Engelfried, 2019))

\[
\text{BR}(K^+ \rightarrow \pi^+\nu\bar{\nu}) < 1.85 \times 10^{-10} \tag{75}
\]
can be used, if we assume the dark photon to decay into light dark-sector fermions, and yields
\[ \frac{\Lambda^2}{\sqrt{\alpha_e d_M^{s_d}}} \gtrsim 9.5 \times 10^6 \text{TeV}^2, \]
which is the strongest among all the limits on the dipole interaction we have discussed.

- Laboratory physics. An interesting limit is derived by means, again, of the data from SN 1987A, this time indirectly from the counting of events in the Kamiokande detector. Axions from the star can, via inverse Bremsstrahlung, excite the oxygen nuclei in the water tank as, in the process \( ^{16}\text{O} \rightarrow ^{16}\text{O}^* \), which subsequently decay producing \( \gamma \) rays triggering the detector. The failure of observing these extra events excludes the values for the coupling \( \alpha'_{aN} \) (Engel et al., 1990)
\[ \quad 6.5 \times 10^{-14} \leq \alpha'_{aN} \leq 8.0 \times 10^{-8}, \]
which can be turned, taking the lower limit in Eq. (77), in
\[ \quad \frac{\Lambda^2}{\sqrt{\alpha_e d_M^{s_d}}} \gtrsim 1.9 \times 10^3 \text{TeV}^2, \]
for the massless dark photons. The limit in Eq. (78) nicely closes the range left open by Eq. (50). A thin windows between Eq. (51) and Eq. (77) is apparently left open for \( \alpha'_{aN} \sim 10^{-14} \).

- Collider physics. Limits from colliders are weaker but are worthwhile to be reported since they come from laboratory physics which is independent of all astrophysical assumptions. The process of pair annihilation into a dark and an ordinary photon provides a striking benchmark (mono-photon plus missing energy) for this search. It applies to electrons in searches at the LEP (Abbiendi et al., 1999; Acciarri et al., 1999; Heister et al., 2003):
\[ \frac{\Lambda^2}{\sqrt{\alpha_e d_M^{s_d}}} \gtrsim 1.2 \text{TeV}^2, \]
and the first generation of quarks at the LHC from CMS (Aaboud et al., 2016) with luminosity of 35.9 fb\(^{-1}\) (the ATLAS result (Sirunyan et al., 2019b) is with smaller luminosity and less stringent):
\[ \quad \frac{\Lambda^2}{\sqrt{\alpha_e d_M^{s_d}}} \gtrsim 4.3 \text{TeV}^2. \]
We computed the limits in Eq. (79) and Eq. (80) for this review by requiring that the number of dark photon events be, bin by bin, less than the difference between the observed and the expected number of events.

3. Can the massless dark photon be seen at all?

The limits for the dark dipole of the massless dark photon, as summarized in Fig. 4 and Fig. 5, are indeed very stringent. For an effective scale \( \Lambda \) around 1 TeV, for example, only values of dipole moments of \( O(10^{-6}) \) for electrons and \( O(10^{-5}) \) for quarks are still allowed. These are numbers making detection in an experiment very challenging.

This does not mean that the massless dark photon cannot be searched for in the laboratory. We must look either to processes where SM particles heavier than the electron or the muon and the \( u \) or \( d \) quarks are involved—and the most severe astrophysical bounds do not apply—or physics where the dipole operator in Eq. (15) is between fermions of different flavors or very high-energy processes where the large scale \( \Lambda \) is partially compensated by the scaling of the dipole and radius operators in Eq. (14) and Eq. (16) and the overall contribution is less suppressed.

For example, for a first generation quark taken to be a parton in a hadron collider, the limit at an energy scale of 10 TeV, is of \( d_M^{s_d} \sim 10^{-3} \) (see Fig. 5) which would give a deviation in the cross section within the reach of future machines. Similarly, for the electron, the limits in Fig. 4 show that a \( d_M^{s_e} \sim 10^{-6} \) is still allowed at the scale of 1 TeV and therefore accessible at future lepton colliders for the projected sensitivity. As much suppressed as these cross sections are, they are comparable with those of the case of the massive dark photon after the corresponding limits are taken into account (see section III.A.1).

These, and others possibilities, are discussed in section II.D where some of the proposed experiments to search for the massless dark photon are reviewed.

B. Limits on milli-charged particles

Milli-charged particles arise, as discussed in the section I.A of the Introduction, in the case of a massless dark photon because the rotation of the mixing term in Eq. (1) leaves the photon coupled to the dark sector particles \( \chi \) with strength \( \varepsilon' \). Searches are accordingly parameterized in terms of the mass \( m_\chi \) and the electromagnetic coupling (modulated by \( \varepsilon' \)) of the supposed milli-charged dark-sector particle.

The physics of stellar evolution for horizontal branches, red giants, and white dwarves (RGWD (Vogel and Redondo, 2014)), together with supernovae (SN1987 (Chang et al., 2018)) provide bounds in the region of small masses \( m_\chi \lesssim 1 \text{ MeV} \). In this region constraints on \( N_{\text{eff}} \) during nucleosynthesis and in the cosmic microwave background (\( N_{\text{eff}} \) BBN and CMB (Vogel and Redondo, 2014)) limit the possibility of having milli-charged particles. These limits are derived along the same lines discussed in the case of the massless dark photon.
FIG. 6. Existing limits (top, filled areas) and future sensitivities for existing or proposed experiments (bottom, colored curves) on milli-charged dark-sector matter. Top plot: Existing limits from stellar evolution (RGWD (Vogel and Redondo, 2014) and SN1987 (Chang et al., 2018)); $N_{\text{eff}}$ during nucleosynthesis and in the cosmic microwave background (Vogel and Redondo, 2014); invisible decays of ortho-positronium (oPS) (Badertscher et al., 2007); SLAC milliQ experiment (Prinz et al., 1998); reinterpretation of data from LSND and miniMooNE (Magill et al., 2019); searches at LEP (Davidson et al., 2000) and LHC (Jaeckel et al., 2013); WMAP results and dark matter relic density abundance (Jaeckel et al., 2013). Bottom plot: Future sensitivities of NA64(e)$^{++}$ (Banerjee et al., 2018a); NA64(µ) (Banerjee et al., 2018b); FerMINI (Kelly and Tsai, 2019); milliQAN (Ball et al., 2016); LDMX (Akesson et al., 2018b). The sensitivity shown for LDMX @ CERN corresponds to $10^{16}$ electrons-on-target and a beam energy of 16 GeV. The existing limits are shown as gray areas. The bottom plot is revised from (Beacham et al., 2020).

Further limits can be derived from precision measurements in QED, notably from the Lamb shift in the transition $2S_{1/2} \rightarrow 2P_{3/2}$ in the Hydrogen atom (Hagley and Pipkin, 1994) and the non-observation of the invisible decay.
of ortho-positronium (oPS (Badertscher et al., 2007)). Limits in the intermediate mass range 1 − 100 MeV come from a SLAC dedicated experiment (SLAC milliQ (Prinz et al., 1998)) and from the reinterpretation of data from the neutrino experiments LSND and miniBooNE (Magill et al., 2019).

Searches at LEP (Davidson et al., 2000) and LHC (Jaeckel et al., 2013) cover larger values of the mass (100 MeV ≲ mχ ≲ 1 TeV).

Finally, for very large masses (mχ ≳ 10 TeV) the impact on the cosmological parameters severely restricts the possible values of milli-quantum charges (WMAP and dark matter relic density constraint, (Jaeckel et al., 2013) and references therein).

All these limits are shown as filled area in the top plot of Fig. 6.

Milli-quantum particles as dark matter have been proposed (see for example (Kovetz et al., 2018) and (Liu et al., 2019)) to explain the anomalous 21 cm hydrogen absorption signal reported by the EDGES experiment (Monsalve et al., 2018). Given the preliminary nature of the results, we have not included them in Fig. 6.

The projected limits of future experiments are depicted in Fig. 6 (bottom) together with the current limits in gray background to show the expected advances. Of these, the most significant for masses around 1 GeV comes from the proposed milliQAN experiment (Ball et al., 2016) proposed to be installed on the surface above one of the LHC interaction points. MilliQAN could improve the collider limits by two orders of magnitude. The range in mass between 10-100 MeV can be optimally covered by the FerMINI experiment (Kelly and Tsai, 2019) proposed in the DUNE near detector hall at Fermilab. Finally the search for milli-charged particles below 10 MeV mass may be improved by almost two orders of magnitude by the LDMX experiment (Alesson et al., 2018b) proposed both at CERN (Alesson et al., 2018a) and at SLAC (Rauenheimer et al., 2018).

C. A minimal model of the dark sector

As discussed in section I, it is useful to underpin the phenomenology of the massless dark photon to a UV model. We consider a minimal model consisting of dark fermions that are, by definition, singlets under the SM gauge interactions. These dark fermions interact with the visible sector through a portal provided by scalar messengers which carry both SM and dark-sector charges. These scalars are phenomenologically akin to the sfermions of supersymmetric models.

In general, we can have as many dark fermions as there are in the SM; they can be classified conveniently according to whether they couple (via the corresponding messengers) to quarks (qL, qR, dR) or leptons (lL, εR): We denote the first (hadron-like) Q and the latter (lepton-like) γ.

The Yukawa-like interaction Lagrangian for flavor-diagonal interactions can be written as (Gabrielli et al., 2017; Gabrielli and Raidal, 2014):

\[
\mathcal{L} \supset -g_L \left( \phi_L^+ \bar{\chi}_L l_L^c + S_{R}^{U,U^r} Q_{R^U} q_L + S_{L}^{D,D^r} Q_{R^D} q_L \right) - g_R \left( \phi_R^+ \bar{\epsilon}_R e_R^c + S_{R}^{U,U^r} Q_{R^U} u_R + S_{R}^{D,D^r} Q_{R^D} d_R \right) + \text{h.c.} \tag{81}
\]

where qL (qR) and εL (εR) are SU(2)L doublets (singlets) for quarks and leptons respectively. Sum over flavor and color indices, that we omitted for simplicity, is understood. The L-type scalars are doublets under SU(2)L, while the R-type scalars are singlets under SU(2)L. The S_{L,R} messengers carry color indices (unmarked in (81)), while the messengers φ_{L,R} are color singlets. The Yukawa coupling strengths are parameterized by α_{L,R} ≡ g_{L,R}^2/(4π); they can be different for different fermions and as many as the SM fermions. For simplicity, we take them to be equal and, in addition, α_L = α_R.

In order to generate chirality-changing processes, we must have the mixing terms

\[
\mathcal{L} \supset -\lambda_S S_0 \left( H^† \phi_R^† \phi_L + \bar{H}^† S_{R}^{U,U^r} S_{L}^{U,U^r} + H^† S_{R}^{D,D^r} S_{L}^{D,D^r} \right) + \text{h.c.}, \tag{82}
\]

where H is the SM Higgs boson, \(\bar{H} = i\sigma_2 H^*\), and S_{0} a scalar singlet of the dark sector. After both S_{0} and H take a vacuum expectation value (\(v_H\) and \(v_h\)—the electroweak vacuum expectation value—respectively), the Lagrangian in Eq. (82) gives rise to the mixing between right- and left-handed states.

Dark sector and messenger states are both charged under an unbroken U(1)D gauge symmetry which is the same of the corresponding massless dark photon, with coupling strength \(\alpha_D\). We assign different dark U(1)D charges to the various dark sector fermions to ensure, by charge conservation, their stability. Since SM fields are neutral under U(1)D interactions, messengers and associated dark-fermions field in Eq.(81) must carry the same U(1)D quantum charge.

When the dark sector scalar S_{0} and the Higgs boson acquire their vacuum expectation values, the scalar messengers must be rotated to identify the physical states. Before this rotation, \(\phi_{L,R}, S_{L,R}, S_{D,R}, S_{L,D}^D \) are degenerate mass eigenstates with mass m_S. After the rotation, the mass eigenstates (labeled by ±) are given by \(\phi_{\pm} \equiv \frac{1}{\sqrt{2}} \left( \phi_L \pm \phi_R \right), S_{L,D}^{U,D} \equiv \frac{1}{\sqrt{2}} \left( S_{L,D}^{U,D} \pm S_{L,D}^{U,D} \right), \) cor-
responding to masses

\[
m_{\pm} = m_{\phi,S} \sqrt{1 \pm \eta_s}
\]

(83)

where we defined the mixing parameters for the S and φ messengers

\[
\lambda_{\phi,S} \equiv \frac{\lambda_{s} M_{\phi,S} v_{h}}{m_{S}^2}.
\]

(84)

In the new basis, the interaction terms in Eq. (81) in the lepton sector is given by

\[
\mathcal{L}^{(lep)} \supset -g_{L} \phi_{L}^{\dagger} (\bar{\chi}_R \nu_{L}) - \frac{g_{L}}{\sqrt{2}} (\phi_{+}^{i} + \phi_{-}^{i} \chi_L e_{R}) - \frac{g_{R}}{\sqrt{2}} (\phi_{-}^{i} - \phi_{+}^{i}) (\bar{\chi}_L \nu_R) + h.c.
\]

(85)

The corresponding interaction terms in the hadronic sector have the same form.

Looking at (85), we can see that if \( \chi \) is a stable dark-sector species, then its mass must be at most \( m_{-} + m_{e} \). Similarly, for a dark-sector species Q, the mass must be no heavier than \( m_{-} + m_{q} \), where \( m_{q} \) is the mass of the SM species corresponding to \( Q \). This sets an upper bound for the mixing \( \eta_{\phi,S} \):

\[
\eta_{\phi,S} < 1 - \left( \frac{M}{m_{\phi,S}} \right)^2.
\]

(86)

In Eq. (86), \( M \) is the mass of the heaviest stable dark-sector species. We assume that \( M \) is heavier than any SM species. The upper bound in Eq. (86) also guarantees that the scalar messengers are heavier than the dark fermion into which they can thus decay.

This model can be considered as a template for many models of the dark sector with the scalar messenger as stand-in for more complicated portals. It is a simplified version of the model in (Gabrielli and Raidal, 2014), which might provide a natural solution to the SM flavor-hierarchy problem.

The discussion above is restricted to the flavor-diagonal interactions. A more general flavor structure in the portal interaction, including the off-diagonal terms, arising as a consequence of the simultaneous diagonalization of the dark-fermion mass and quark interaction basis, can be simply obtained by generalizing the above terms as follows (Gabrielli et al., 2016)

\[
S_{L}^{11} Q_{R}^{1} q_{L}^{1} \rightarrow S_{L}^{11} Q_{R}^{1} (\rho_{L})^{11} q_{L}^{1},
\]

(87)

\[
S_{L}^{i j} Q_{R}^{i} q_{L}^{j} \rightarrow S_{L}^{i j} Q_{R}^{i} (\rho_{L})^{i j} q_{L}^{j},
\]

and analogously for the down and lepton sectors, where \( i, j \) are explicit flavor indices and sum over \( i, j \) is understood.

To keep the contribution to the dipole coefficient simple, lest the generality obfuscates the estimate, we follow the guidelines of the model in (Gabrielli and Raidal, 2014). We assume that the masses of the messengers \( \phi^{1}, S_{L,R}^{1}, \text{and } S_{L,R}^{2} \) are the same and the mixing matrices \( \rho_{ij} \) have a hierarchical structure (like in the SM) with the off-diagonal smaller than the diagonal terms. The former hypothesis is a consequence of the \( SU(N_{F}) \) flavor symmetry in the free lagrangian of messenger sector (with \( N_{F} = 6 \)) (Gabrielli and Raidal, 2014), while the latter follows from the requirement of minimal flavor violation hypothesis (Gabrielli et al., 2017).

We also take \( \rho_{ij} \equiv \rho_{ij}^{D} = \rho_{ij}^{S} \). This way, the loop of dark sector particles is dominated by the contribution with the heaviest dark fermion coupled to the SM fermions of flavor \( i \) and \( j \) with one coefficient off-diagonal \( \rho_{ij} \) and one diagonal \( \rho_{ii} \). In the following, in order to distinguish the contribution from the up and down sector couplings we will use the notation \( \rho_{uu} \equiv \rho_{11}^{U} \), \( \rho_{dd} \equiv \rho_{11}^{D} \), \( \rho_{sd} \equiv \rho_{21}^{D} \), and similarly for the other coefficients.

Matching the model to the effective Lagrangian given in Eq. (14) after integrating the loop, and identifying the scale \( \Lambda \) as

\[
\frac{v_{h}}{\Lambda^{2}} \simeq \frac{m_{Q}^{2}}{m_{S}^{2}},
\]

(88)

with \( m_{Q} \), the heaviest dark-fermion running in the loop, we can re-express the magnetic dipole explicitly in terms of the parameters of the model. For example, in the case of the generic (quark) flavor transition from \( i \rightarrow j \), with \( i \)- and \( j \)-mixing, neglecting the SM masses, according to the Lagrangian in (81) and substitutions (87), we have (Gabrielli et al., 2016)

\[
D_{M}^{ij} = \rho_{ij} \rho_{ij}^{*} \frac{g_{L} g_{R}}{(4\pi)^{2}} F_{M}(x, \eta_{s}).
\]

(89)

where \( x = (m_{Q})^{2}/m_{S}^{2} \) and \( \eta_{s} \) the mixing parameter defined in (84). In the following, we will introduce the
notation of $m_{S\ell}$ and $m_{S\ell}$ to distinguish the common messenger mass in the up and down $SU(2)_L$ sectors respectively, and $\eta_{\alpha,\beta}$ for the corresponding mixing parameters. The function $F_M(x, y)$ is given by (Gabrielli et al., 2016)

$$F_M(x, y) = \frac{1}{2} \left[ f(x, y) - f(x, -y) \right],$$

(90)

where

$$f(x, y) = \frac{1 - x + y + (1 + y) \log \left( \frac{x}{1 + y} \right)}{(1 - x + y)^2}.$$  

(91)

CP-violating phases, relevant for flavor changing processes, can arise from the mixing parameters. For instance, in the $n \rightarrow m$ flavor transition, we can have CP-violating phase $\delta_{CP}$ from the relation

$$\rho_{nm}\rho_{mm}^* - \rho_{nm}^*\rho_{mm} = 2i \sin \delta_{CP}.$$  

(92)

1. Constraints on the UV model parameters

The introduction of the UV model makes possible to re-discuss the bounds of section II.A on the massless dark photon in terms of the parameters of the model.

There are no laboratory limits for the masses of the dark fermions from events in which they are produced because they are SM singlets and do not interact directly with the detector. Cosmological bounds have been considered in (Acuña et al., 2020) where, in particular, avoiding distortions of the cosmic microwave background is shown to require the masses of the dark fermions to be larger than 1 GeV or, if lighter, that the coupling $\alpha_L$ and $\alpha_R$ be less than $10^{-3}$.

The messenger states have the same quantum numbers and spin as the supersymmetric squarks. At the LHC they are copiously produced in pairs through QCD interactions and decay at tree level into a quark and a dark fermion. The final state arising from their decay is thus the same as the one obtained from the $\tilde{g} \rightarrow q\chi_0^0$ process. Therefore limits on the messenger masses can be obtained by reinterpreting supersymmetric searches on first and second generation squarks decaying into a light jet and a massless neutralino (Aaboud et al., 2018), assuming that the gluino is decoupled. A lower bound on their masses is thus obtained (Barducci et al., 2018) to give

$$m_S^i \gtrsim 940 \text{ GeV},$$

(93)

for the messenger mass related to the dark fermions $Q^U$ and $Q^D$. This limit increases up to 1.5 TeV by assuming that messengers of both chiralities associated to the first and second generation of SM quarks are degenerate in mass.

For the masses of the lepton-like scalar messengers, constraints on the mass of sleptons (Sirunyan et al., 2019c) give the following lower bound on the messenger mass in the lepton sector:

$$m_{\phi} \gtrsim 290 \text{ GeV}.$$  

(94)

All the limits discussed in section II.A can be re-expressed in terms of the UV model parameters.

For example, the limit from stellar cooling in Eq. (44) becomes

$$\frac{m_{\phi}^2/m_{\chi^\pm}}{\sqrt{\alpha_D\alpha_L\alpha_R} |\rho_{ee}|^2} F_M(x_e, \eta_\phi) \gtrsim 2.1 \times 10^6 \text{ TeV},$$  

(95)

where $x_e = m_{\phi}^2/m_{\chi^\pm}^2$, with $m_{\chi^\pm}$ the dark-fermion mass associated to the electron, and $\eta_\phi$ the corresponding mixing parameter in the colorless messengers sector, and the loop function $F_M(x, y)$ is given in Eq. (90). This limit, which is obtained by rescaling the right-hand side of Eq. (44) for $1/(4\pi v_h)$, applies specifically to the Yukawa coupling of electrons and the corresponding messenger state.

For a quick estimate of the bound above and those that follow, the loop function $F_M(x, y)$ can be considered a coefficient of order $O(10^{-1})$ as long as $\eta_\phi$ is not too small. For instance, for $x \approx 1$ and $y \approx 0.5$, the loop function $F_M \approx 0.09$.

Similarly, by using the same rescaling factor, the neutrino signal of supernova 1987A and the limit in Eq. (50) yields now

$$\frac{m_{S}\bar{m}_{Q_e}}{\sqrt{\alpha_D\alpha_L\alpha_R} |\rho_{uu}|^2} F_M(x_u, \eta_S) \gtrsim 2.0 \times 10^5 \text{ TeV},$$  

(96)

where now $x_u = m_{S}^2/m_{Q_e}^2$, with $m_{Q_e}$ the dark-fermion associated to the light $u$ quark. A similar limit holds for the case of the $d$ quark sector.

The others bounds in section II.A can be written in terms of the parameters of the model in the same way.

Instead, new bounds can be set now that we have an underlying UV model because the scalar messengers carry also the electromagnetic charge. Processes with the visible photon can thus be used; these processes were not available for the model-independent case in section II.A for which only the coupling to the dark photon was taken into account.

The magnetic moment of the SM fermions arises from the one-loop diagram of the states of the UV model. From Eq. (68) in section II.A, we find

$$\frac{m_{\phi}^2/m_{\chi^\pm}}{\sqrt{\alpha_D\alpha_L\alpha_R} |\rho_{ee}|^2} G_M(x_e, \eta_\phi) \gtrsim 9.8 \times 10^4 \text{ TeV},$$  

(97)

where $x_e = m_{\phi}^2/m_{\chi^\pm}^2$, with $m_{\chi^\pm}$ the dark-fermion mass associated to the muon. The loop function is in this case given by (Gabrielli et al., 2016)

$$G_M(x, y) = \frac{1}{2} \left[ g(x, y) - g(x, -y) \right],$$  

(98)
where
g(x, y) = \frac{(1 + y)^2 - x^2 + 2x(1 + y) \log \left( \frac{x + y}{1 + y} \right)}{2(x - 1 - y)^3}. \quad (99)

Also interesting is the anomalous magnetic moment of the muon because of the lingering discrepancy between theory and experiments. From Eq. (70) in section II.A, we find

\[
m^2_{\phi}/m_{\chi}^2 \gtrsim 6.3 \times 10^3 \text{ TeV},
\]

where \(x_\mu = m^2_{\chi}/m^2_{\phi}\), with \(m_{\chi}\) the dark-fermion mass associated to the muon. Again, for a quick estimate of the bounds above and those that follow, the loop function \(G_M(x, y)\) can be considered a coefficient of order \(O(10^{-1})\) as long as \(\eta_0\) is not too small. For instance, for \(x \approx 1\) and \(y \approx 0.5\), the loop function \(G_M \approx 0.05\).

The various Yukawa couplings and messenger and fermion masses are probed in a selective manner in flavor physics where we must distinguish among the various couplings and states. Mixing (proportional to a coefficient \(p_{ij}\) in the equations below) between different flavor states must be included.

The strongest bound comes from the limit on the BR(\(\mu \to e\gamma\)) < 4.2 \times 10^{-13} (CL 90%) (Baldini et al., 2016) of the MEG experiment. From this result, we find that

\[
m^2_{\phi}/m_{\chi}^2 \gtrsim 6.3 \times 10^3 \text{ TeV},
\]

where \(x_\mu = m^2_{\chi}/m^2_{\phi}\), with \(m_{\chi}\) the dark-fermion mass associated to the muon. Again, for a quick estimate of the bounds above and those that follow, the loop function \(G_M(x, y)\) can be considered a coefficient of order \(O(10^{-1})\) as long as \(\eta_0\) is not too small. For instance, for \(x \approx 1\) and \(y \approx 0.5\), the loop function \(G_M \approx 0.05\).

A weaker bound can be extracted, in the hadronic sector, from the difference between the experimental limit on the BR(\(B \to X_s\gamma\)) < (3.21 ± 0.33) \times 10^{-4} (Lees et al., 2012b) of the BaBar collaboration and its SM estimate (Misiak et al., 2007). It yields

\[
m^2_{\phi}/m_{\chi}^2 \gtrsim 4.9 \times 10^8 \text{ TeV}.
\]

The limits in Eq. (101) and Eq. (102) apply specifically to the off-diagonal terms in the Yukawa couplings \(p_{ij}\) of the muon-electron and \(b-s\) quark mixing, respectively, and to the corresponding mass of messenger states.

The mass mixing in the Kaon system (Fabbrichesi et al., 2017; Gabrielli et al., 2016) gives a further limit

\[
m^2_{\phi}/m_{\chi}^2 \gtrsim 1.3 \times 10^4 \text{ TeV},
\]

where \(x_b = m^2_{\chi}/m^2_{\phi}\), with \(m_{\chi}\) the mass of dark fermion associated to the \(b\)-quark.

The limits in Eq. (101) and Eq. (102) apply specifically to the off-diagonal terms in the Yukawa couplings \(p_{ij}\) of the muon-electron and \(b\)-\(s\) quark mixing, respectively, and to the corresponding mass of messenger states.

The mass mixing in the Kaon system (Fabbrichesi et al., 2017; Gabrielli et al., 2016) gives a further limit

\[
m^2_{\phi}/m_{\chi}^2 \gtrsim 3 \times 10^5 \text{ TeV}^2,
\]

which is not related to the dark photon and its coupling \(\alpha_d\) because it comes from the box-diagram insertion of the dark scalars and fermions.

The limit in Eq. (103) is obtained by requiring that the messenger contribution to the box diagram for the \(K^0\) shows that \(\eta_0\) does not exceed the experimental value of the mixing parameter \(\Delta m_K = 3.48 \times 10^{-12}\text{MeV}\) (Tanabashi et al., 2018). Due to chirality arguments, the leading contribution to the box diagram in Eq. (103) does not depend on the dark fermion mass, which is assumed to be much smaller than the corresponding messenger mass in the down sector and therefore very weakly on the loop function.

The limit in Eq. (103) applies specifically to the off-diagonal term in the Yukawa coupling of \(d-s\) quark mixing and the corresponding messenger state. A similar but weaker bound can be found from \(B\)-meson mixing.

As displayed in the equations above, all these limits can be made weaker by taking \(m_{\chi}\) (or \(m_{\phi}\)) sufficiently light or by varying the corresponding mixing parameters \(\eta_s\). In the UV model is thus possible to play with the parameters to make room for larger values of the dipole coefficient by absorbing part of the suppression in the connection between the scale \(\Lambda\) and the mass ratios \(m_{\chi}/m_{\phi}^2\) and \(m_{\chi}/m_{\phi}\). For instance a scale \(\Lambda = 1\text{ TeV}\) for the new physics of the dark sector is still allowed by the stringent bound in Eq. (70) if we take \(m_{\chi}\) sufficiently small. This way, there is some additional freedom in comparing limits from different processes as compared to the model-independent case where the scale \(\Lambda\) is taken to be the same for all bounds.

D. Future experiments

The massless dark photon has been neglected so far from the experimental point of view as compared to the massive one. It is one of the aims of the present review to boost the community scrutiny in this direction. In the past few years several proposals have been put forward and new experiments are in the planning:

- Flavor physics: This is one of the most promising areas for searching for the dark photon and the dark sector in general because none of the stringent astrophysical constrains discussed in section II.A applies given the flavor off-diagonal nature of the dipole operator in these cases.

Proposals exist for processes in Kaon physics at NA62 (Cortina Gil et al., 2017). The Kaon decay \(K \to \pi A'\) is forbidden by the conservation of angular momentum but the decay \(K^+ \to \pi^0 \pi^+ A'\) is allowed and the estimated branching ratio (Fabbrichesi et al., 2017) is within reach of the current sensitivity. The rare decays \(K^+ \to \pi^+ \nu \bar{\nu}\) (Cortina Gil et al., 2019b) and \(K_{L} \to \pi^0 \nu \bar{\nu}\) (Ahn et al., 2019) are other two processes where the physics of the dark photon can play a crucial role (Fabbrichesi and Gabrielli, 2019). Also Hyperion decays can be used
for detecting the production of $A'$ (Su and Tandean, 2019) and in the decay of charmed hadrons (Su and Tandean, 2020) and BESIII.

In addition, decays into invisible states of $B$-mesons at BaBar (Lees et al., 2012a) and Belle (Hsu et al., 2012) and $K_{L,S}$ and other neutral mesons at NA64 (Gninenko, 2015; Gninenko and Krasnikov, 2015) can be used to study the dark sector (assuming the invisible states belong to it). These decays are greatly enhanced by the Fermi-Sommerfeld (Fermi, 1934; Sommerfeld, 1931) effect due to their interaction with the dark photon—the same way as ordinary decays, like the $\beta$-decay, are enhanced by the same effect—making this another exciting area for searching the dark sector (Barducci et al., 2018).

- **Higgs and Z physics:** The striking signature of a mono-photon plus missing energy can be used to search Higgs (Biswas et al., 2015, 2017; Gabrielli et al., 2014) and Z-boson (Fabbrichesi et al., 2018) decay into a visible and a dark photon. Again, the stringent astrophysical constrains discussed in section II.A do not apply because the size of the dipole operator is dominated (in the loop diagram) by the heavy-quark contribution’s giving raise to the coupling to the dark photon, as discussed in section II.C.

- **Pair annihilation:** Collider experiment at higher energies and luminosities can use the same striking signature of a mono-photon plus missing energy to search for the dark photon. Even though the dipole interaction is suppressed and severely constrained in this case by the astrophysical and cosmological bounds discussed in section II.A, it is no more suppressed than the equivalent cross sections for the massive case. Moreover, the dipole operator scales as the center-of-mass energy in the process and higher energies make it more and more relevant;

- **Magnons:** An interesting possibility is the use of magnons in ferromagnetic materials and their interaction with dark photons (QUAX proposal) (Barbieri et al., 2017; Chigusa et al., 2020). The estimated sensitivity is again done for axions but can be translated for massless dark photons as in the discussion about stars above.

- **Astrophysics:** Gravitation waves emitted during the inspiral phase of neutron star collapse can test the presence of other forces beside gravitation. Dipole radiation by even small amount of charges on the stars modifies the energy emitted; the dark photon is a prime candidate for this kind of correction (Alexander et al., 2018; Croon et al., 2018; Fabbrichesi and Urbano, 2019; Kopp et al., 2018).

III. PHENOMENOLOGY OF THE MASSIVE DARK PHOTON

The phenomenology of the massive dark photon is discussed in terms of its interaction with the SM particles, as given by Eq. (5):

$$\mathcal{L} = \varepsilon e J^\mu A'_\mu,$$

where $J^\mu$ is the electromagnetic current. The strength of this interaction is modulated by the parameter $\varepsilon$. The parameter space for the experimental searches is given by the mass of the dark photon $m_{A'}$ and the mixing parameter $\varepsilon$.

1. Production, decays and detection

Because the current in Eq. (104) is the same as the usual electromagnetic current, dark photons $A'$ can be produced like ordinary photons. The main production mechanisms are:

- **Bremsstrahlung.** The incoming electron scatters off the target nuclei ($Z$), goes off-shell and can thus emit the dark photon: $e^- Z \rightarrow e^- Z A'$;

- **Annihilation:** An electron-positron pair annihilates into an ordinary and a dark photon: $e^- e^+ \rightarrow \gamma A'$;

- **Meson decays:** A meson $M$ (it being a $\pi^0$, $\eta$, or a $K$ or a $D$) decays as $M \rightarrow \gamma A'$;

- **Drell-Yan:** A quark-antiquark pair annihilates into the dark photon, which then decays into a lepton pair (or hadrons): $q\bar q \rightarrow A'(\rightarrow \ell^+\ell^-$ or $h^+h^-)$.
Different experiments use different production mechanisms and, sometime, more than one simultaneously.

Detection of $A'$ is based on its decays modes. The decay width of the massive dark photon $A'$ into SM leptons $\ell$ is

$$\Gamma(A' \to \ell^+ \ell^-) = \frac{1}{3} \alpha \varepsilon^2 m_{A'} \sqrt{1 - \frac{4m_e^2}{m_{A'}^2}} \left(1 + \frac{2m_e^2}{m_{A'}^2}\right),$$

which is only open for $m_{A'} > 2m_e$. Similarly, the width into hadrons is

$$\Gamma(A' \to \text{hadrons}) = \frac{1}{3} \alpha \varepsilon^2 m_{A'} \sqrt{1 - \frac{4m_e^2}{m_{A'}^2}} \left(1 + \frac{2m_e^2}{m_{A'}^2}\right) R,$$

where $R \equiv \sigma_{e^+e^-\to \text{had}}/\sigma_{e^+e^-\to \mu^+\mu^-}$.

Since all visible widths are proportional to $\varepsilon$, the branching ratios are independent of it.

At accelerator-based experiments, several approaches can be pursued to search for dark photons depending on the characteristics of the available beam line and the detector. These can be summarized as follows:

- Detection of visible final states: dark photons with masses above $\sim 1$ MeV can decay to visible final states. The detection of visible final state is a technique mostly used in beam-dump and collider experiments, where typical signatures are expected to show up as narrow resonances over an irreducible background. Collider experiments are typically sensitive to larger values of $\varepsilon$ ($\varepsilon > 10^{-3}$) than beam dump experiments which typically cover couplings below $10^{-3}$.

  Missing momentum/energy techniques: invisible decay of dark photons can be detected in fixed-target reactions as, for example, $e^- Z \to e^- Z A'$ ($Z$ being the nuclei atomic number) with $A' \to \chi \gamma$ and $\chi$ being a putative dark matter particle, by measuring the missing momentum or missing energy carried away from the escaping invisible particle or particles. The main challenge for this approach is the very high background rejection that must be achieved, which relies heavily on the detector hermeticity and, in some cases, on the exact knowledge of the initial and final state kinematics.

- Missing mass technique: this technique is mostly used to detect invisible particles (as dark-matter candidates) in reactions with a well-known initial state, as for example $e^+e^- \to A' \gamma$, with $A' \to \chi \chi$. It requires detectors with very good hermeticity that allow to detect all the other particles in the final state. Characteristic signature of this reaction is the presence of a narrow resonance emerging over a smooth background in the distribution of the missing mass. The main limitation of this technique is the required knowledge of the background arising from processes in which particles in the final state escape the apparatus without being detected.

2. Visible and invisible massive dark photon

In collecting the limits on the parameters of massive dark photon is important to distinguish two cases accordingly on whether its mass is smaller or larger than twice the mass of the electron, the lightest charged SM fermion.

The dark photon is visible if its mass is $M_{A'} > 2m_e \sim 1$ MeV because it can decay into SM charged states which leave a signature in the detectors. We discuss the limits on the visible dark photon in section III.A.1.

In the same regime for which $M_{A'} > 1$ MeV, however, the massive dark photon could also decay into dark sector states if their masses are light enough. In this case we have a non-vanishing branching ratio into invisible final states. The invisible decay into these states of the dark sector $\chi$ in given by

$$\Gamma(A' \to \chi \bar{\chi}) = \frac{1}{3} \alpha_D m_{A'} \sqrt{1 - \frac{4m_e^2}{m_{A'}^2}} \left(1 + \frac{2m_e^2}{m_{A'}^2}\right).$$

Dark photons decays into this invisible channel if $m_{A'} > 2m_e$; this channel dominates if $\alpha_D \gg \alpha\varepsilon^2$.

Most of the experimental searches with dark photon in visible decays assume that the dark-sector states are not kinematically accessible and the dark photon is visible only through its decay into SM states. The limits need to be re-modulated if the branching ratio into invisible states is numerically significant or even dominant. We discuss this case in section III.A.2 below.

FIG. 9 Decay of the massive dark photon into visible (SM leptons or hadrons) and invisible (DM) modes
If the mass of the dark photon is less than 1 MeV, it cannot decay in any known SM charged fermion and its decay is therefore completely invisible. The experimental searches for dark photon into invisible final states are based on the energy losses that the production of dark photons, independently of his being stable or decaying into dark fermions, implies on astrophysical objects like stars or in signals released in direct detection dark matter experiments. The experimental limits in the case of the invisible dark photon are discussed in section III.A.3 below.

A. Limits on the parameters $\varepsilon$ and $m_{A'}$

As discussed, the space of the parameters (the mixing $\varepsilon$ and the mass $m_{A'}$ of the dark photon) is best spanned in two regions according on whether the mass $m_{A'}$ is larger or smaller than twice the mass of the electron: Roughly 1 MeV.

1. Constraints for $m_{A'} > 1$ MeV with $A'$ decays to visible final states

Two kinds of experiments provide the existing limits on the visible massive dark photon in the region of $m_{A'} > 1$ MeV: experiments at colliders and at fixed-target or beam dumps. In both cases the experiments search for resonances over a smooth background, with a vertex prompt or slightly displaced with respect to the beam interaction point in case of collider, or highly displaced in case of beam dump based experiments. The two categories are highly complementary, being the first category mostly sensitive to relatively large values of the mixing parameter $\varepsilon$, ($\varepsilon > 10^{-3}$) and the dark photon mass (up to several tens of GeV for $pp$ collider experiments), while the second is sensitive to relatively small values ($10^{-7} \lesssim \varepsilon \lesssim 10^{-3}$) in the low mass range, $m_{A'}$ less than few GeV.

- Experiments at colliders. These experiments search for resonances in the invariant mass distribution of $e^+e^-$, $\mu^+\mu^-$ pairs. Different dark-photon production mechanisms are used in the different experiments: meson decays ($\pi^0 \rightarrow \gamma A'$, NA48/2 (Batley et al., 2015)), Bremsstrahlung ($e^-Z \rightarrow e^-ZA'$, A1 (Merkel et al., 2014)), annihilation ($e^+e^- \rightarrow \gamma A'$, BaBar (Lees et al., 2014)), and all these processes in different searches at KLOE (Anastasi et al., 2016; Archilli et al., 2012; Babusci et al., 2013, 2014). In a proton-proton ($pp$) collider the dark photon is produced via the $\gamma - A'$ mixing in all the processes where an off-shell photon $\gamma^*$ with mass $m(\gamma^*)$ is produced: meson decays, Bremsstrahlung, and Drell-Yan production. LHCb (Aaij et al., 2018b, 2020) has performed a search for dark photon decaying in $\mu^+\mu^-$ final states using 1.6 fb$^{-1}$ of data collected at the LHC $pp$ collisions at 13 TeV centre-of-mass energy. CMS (Sirunyan et al., 2019a) has performed the same search using 137 fb$^{-1}$ of fully reconstructed data and 96.6 fb$^{-1}$ of data collected with a reduced trigger information.

Fig. 10 (top) shows the limits from experiments at extracted electron beams (E141 (Riordan et al., 1987) and E137 (Batell et al., 2014; Bjorken et al., 1988; Marsiglio et al., 2018) at SLAC, E774 (Bross et al., 1991) at Fermilab) and at extracted proton beams from CHARM at CERN ((Gninenko, 2012) based on CHARM data (Bergsma et al., 1985)).

In addition, bounds on energy losses in supernovae (Dent et al., 2012; Dreiner et al., 2014), as discussed in the massless case, provide further limits in the region of small masses. Also the electron magnetic moment, with its very precise experimental determination, can be used to set an indirect limit (Pospelov, 2009). These limits are included in Fig. 10.

Recent constraints from ATLAS (And et al., 2014, 2016) and CMS (Khachatryan et al., 2016) would nominally cover the interesting region around 1 GeV for $\varepsilon$ between $10^{-6}$ and $10^{-2}$ but unfortunately they have been framed within a restrictive model and are not on the same footing that the limits included in Fig. 10.

Additional limits (not included in Fig. 10) from cosmology (in the cosmic microwave background and nucleosynthesis) exist in the very dark region of very small $\varepsilon < 10^{-10}$ (Fradette et al., 2014).

Looking at Fig. 10 (top) it is clear that it would be desirable to first close the gap between the beam-dump and the collider based experiments in the region between tens of MeV up to 1 GeV in the dark photon mass, and then extend the limits for larger masses. Both of these goals could be achieved through a series of experiments summarized here below whose sensitivity is shown in Figure 10 (bottom) as colored curves.
FIG. 10 Top plot: Existing limits on massive dark photon for $m_{A'} > 1$ MeV from di-lepton searches at experiments at collider/fixed target (A1 (Merkel et al., 2014), LHCb (Aaij et al., 2020), CMS (Sirunyan et al., 2019a), BaBar (Less et al., 2014), KLOE (Anastasi et al., 2016; Archilli et al., 2012; Babusci et al., 2013, 2014), and NA48/2 (Batley et al., 2015)) and old beam dump: E774 (Bross et al., 1991), E141 (Riordan et al., 1987), E137 (Batell et al., 2014; Bjorken et al., 1988; Marsicano et al., 2016), ν-Cal (Blumlein and Brunner, 2011, 2014), and CHARM (from (Gninenko, 2012). Bounds from supernovae (Dent et al., 2012; Dreiner et al., 2014) and $(g - 2)_e$ (Pospelov, 2009) are also included. Bottom plot: Colored curves are projections for existing and proposed experiments: Belle-II (Altmannshofer et al., 2019) at SuperKEKB, LHCb upgrade (Ilten et al., 2016, 2015) at the LHC, NA62 in dump mode (Cortina Gil et al., 2019a) and NA64(e)$^{++}$ (Banerjee et al., 2018a) at the SPS; FASER and FASER2 (Feng et al., 2018) at the LHC; SeaQuest (Berlin et al., 2018) at Fermilab; HPS (Adrian et al., 2018) at JLAB; an NA64-like experiment at AWAKE (Caldwell et al., 2018), and an experiment dedicated to dark photon searches at MESA (Doria et al., 2018, 2019). For masses above 10 GeV projections obtained for ATLAS/CMS during the high luminosity phase of the LHC (HL-LHC (Curtin et al., 2015)) and for experiments running at a future FCC-ee (Karliner et al., 2015), LHeC/FCC-eh (D’Onofrio et al., 2020), and FCC-hh (Curtin et al., 2015) are also shown. The vertical red line shows the allowed range of couplings of a new gauge boson $X$ to electrons that could explain the $^8$Be anomaly (Feng et al., 2016, 2017). The existing limits are shown as gray areas. The bottom plot is revised from (Beacham et al., 2020).

- Belle-II at SuperKEKB will search for visible dark photon decays $A' \rightarrow e^+e^-, \mu^+\mu^-$ where $A'$ is pro-
duced in the process $e^+e^- \rightarrow A'\gamma$. The projections shown in Fig. 10 is based on 50 ab$^{-1}$ of integrated luminosity (Altmannshofer et al., 2019).

- **LHCb upgrade (phase I and phase II) at the LHC**: LHCb phase I will search for dark photon in visible final states both using the inclusive di-muon production (Ilten et al., 2016) and the $D^{*0} \rightarrow D^{0}e^+e^-$ decays (Ilten et al., 2015). The projections are based on 15 fb$^{-1}$, 3 years data taking with 5 fb$^{-1}$/year with an upgraded detector after the LHC Long Shutdown 2. This can be further improved with a possible Phase II upgraded detector (Aaij et al., 2018a) collecting up to 300 fb$^{-1}$ of integrated luminosity after Long Shutdown 4.

- **NA62** or **NA62 in dump mode at the SPS, CERN**: will search for a multitude of feebly-interacting particles, including dark photon, decaying into visible final states and possibly emerging from the interactions of 400 GeV proton beam with a dump. NA62 aims to collect approximately 10$^{18}$ protons-on-target in 2021-2024 (Cortina Gil et al., 2019a).

- **NA64**(e)$^{++}$ at the SPS, CERN, is the upgrade of the existing NA64(e) experiment. It aims to collect about 5 × 10$^{12}$ electrons-on-target after the CERN Long Shutdown 2 (Banerjee et al., 2018a).

- **NA64-like experiment at AWAKE, CERN**: progress in the coming years in proton-driven plasma wakefield acceleration of electrons at the AWAKE facility at CERN could allow an NA64-like experiment to be served by a high-intensity high energy primary electron beam for search for dark photons in visible final states (Caldwell et al., 2018). The sensitivity plot has been obtained assuming ~ 10$^{10}$ electrons-on-target with an energy of 50 GeV.

- **FASER and FASER2 at the LHC, CERN**: FASER (Feng et al., 2018) is being installed in a service tunnel of the LHC located along the beam collision axis, 480 m downstream from the ATLAS interaction point. At this location, FASER (and possibly its larger successor FASER2) will enhance the LHC discovery potential by providing sensitivity to dark photons, dark Higgs bosons, heavy neutral leptons, axion-like particles, and many other proposed feebly-interacting particles (Ariga et al., 2019). FASER and FASER2 aim to collect 150 fb$^{-1}$ and 3000 fb$^{-1}$ of integrated luminosity, respectively.

- **HPS at Jefferson Laboratory (JLab)** The HPS experiment (Adrian et al., 2018), proposed at an electron beam-dump at CEBAF electron beam (2.2-6.6 GeV, up to 500 nA), search for visible ($A' \rightarrow e^+e^-$) dark photon (prompt and displaced) decays produced via Bremsstrahlung production in a thin W target. The experiment makes use of the 200 nA electron beam available in Hall-B at Jefferson Lab.

- **SeaQuest at Fermilab (FNAL)** will search for visible dark photon decays $A' \rightarrow e^+e^-$ at the 120 GeV main injector proton beamline at FNAL (Berlin et al., 2018). It plans to accumulate approximately 10$^{18}$ protons-on-target by 2024.

- **MAGIX or Beam Dump Experiment at MESA, Mainz**: The MESA accelerator is a continuous wave linac that will be able to provide an electron beam of $E_{\text{max}} = 155$ MeV energy and up to 1 mA current (Doria et al., 2019). The MAGIX detector is a twin arm dipole spectrometer placed around a gas target and will search for search for visible ($A' \rightarrow e^+e^-$) dark photon (prompt and displaced) decays produced via Bremsstrahlung production (Doria et al., 2018). The possibility of a beam dump setup experiment is also under study.

**Timeline**: targeted operations in 2021-2022 and 2 years of data taking.

- **Experiments at a future $e^+e^-$ circular collider, FCC-ee**: a powerful technique to be exploited at experiments running at a future $e^+e^-$ circular collider is the radiative return, $e^+e^- \rightarrow A'\gamma, A' \rightarrow \mu^+\mu^-$. The results obtained in (Karliner et al., 2015) have been rescaled to the integrated luminosities of 150 fb$^{-1}$ at $\sqrt{s} = 90$ GeV and 5 ab$^{-1}$ at $\sqrt{s} = 250$ GeV, as in (Ellis et al., 2019).

- **ATLAS/CMS at the high-luminosity phase at the LHC at a future $pp$ circular collider**: At the LHeC (FCC-eh) a 7 TeV (50 TeV) a proton beam collides with a 60 GeV electron beam achieving a center-of-mass energy of 1.3 TeV (3.5 TeV) and a total integrate luminosity of 1 ab$^{-1}$ (3 ab$^{-1}$). At eh colliders the main production process for the dark photon is the deep inelastic scattering $e^-+\text{parton} \rightarrow e^-+\text{parton} A'$, with $A' \rightarrow$ charged fermions (D’Onofrio et al., 2020).
FIG. 11 Existing limits (top, filled areas) and future sensitivities for existing or proposed experiments (bottom, colored curves) for a massive dark photon going to invisible final states ($\alpha_D \gg \alpha_e^2$). Top plot: Existing limits from Kaon decay experiments (E787 (Adler et al., 2002), E949 (Artamonov et al., 2009), NA62 (Cortina Gil et al., 2019c)), BaBar (Lees et al., 2017), and NA64(e) (Banerjee et al., 2019). The constraints from $(g - 2)_\mu$ (Bennett et al., 2006) and $(g - 2)_e$ are also shown. Bottom plot: Future sensitivities for NA64(e)$^{++}$ (Banerjee et al., 2018a), Belle II (Altmannshofer et al., 2019), KLEVER (Ambrosino et al., 2019), PADME (Raggi et al., 2015), LDMX@SLAC (Akesson et al., 2018a,b), and LDMX@CERN (Akesson et al., 2018b; Raubenheimer et al., 2018). The sensitivity curves for LDMX@SLAC and LDMX@CERN assume $10^{14}$ electrons-on-target and $E_{\text{beam}} = 4$ GeV and $10^{16}$ electrons-on-target and $E_{\text{beam}} = 16$ GeV, respectively. The bottom plot is revised from (Beacham et al., 2020). See text for details.

2. Constraints for $m_{A'} > 1$ MeV with $A'$ decays to invisible final states

Different constraints apply in the case of massive dark photon going into invisible final states in the mass region $m_{A'} > 1$ MeV. In this case techniques like missing momentum, missing energy, and missing mass are used in order to identify a possible massive dark photon decaying...
into invisible final states.

The most stringent bounds come from BaBar (Lees et al., 2017) and the electron beam dump NA64(e) experiment at CERN (Banerjee et al., 2019) which recently superseded the results from Kaon experiments (E787 (Adler et al., 2002) and E949 (Artamonov et al., 2009) at BNL, NA62 (Cortina Gil et al., 2019c) at CERN). The existing bounds are depicted in the top plot of Fig. 11 as colored areas. These limits overlap with the exclusion regions defined by the dark photon in the region $3 < m_A' < 10$ MeV and complement them in the range of masses $10$ MeV $< m_A' < 1$ GeV and kinetic mixing strength $10^{-5} < \varepsilon < 10^{-3}$, where the searches of dark photon into visible decays are typically weaker.

Sensitivities of existing or proposed experiments are shown in the bottom plot of Fig. 11 as colored lines. These include:

- $\text{NA64}(e)^{++}$ with $5 \times 10^{12}$ electrons-on-target will search $A' \rightarrow \text{invisible}$ final states with a missing energy technique using a secondary electron beam of $\sim 100$ GeV at the CERN SPS (Banerjee et al., 2018a).

- $\text{Belle II}$ will search for dark photons in the process $e^+e^- \rightarrow A'$ and $A' \rightarrow \text{invisible}$ (Altmannshofer et al., 2019). Projections are based on $20$ fb$^{-1}$ of integrated luminosity.

- $\text{KLEVER}$, proposed at the SPS, could search for dark photons in invisible final states as a by-product of the analysis of the $K_L \rightarrow \pi^0\nu\bar{\nu}$ rare decay, pushing further the investigation performed by traditional Kaon experiments in the mass region between $100-200$ MeV (Ambrosino et al., 2019).

- $\text{PADME}$ (Raggi et al., 2015) will search for $A' \rightarrow \text{invisible}$ final states using the missing momentum technique at the Beam Test Facility (BTF) at Laboratori Nazionali di Frascati (INFN). It will use a $550$ MeV positron beam on a diamond target. A first commissioning run was performed in late 2018 and early 2019 to assess the detector performance and beam line quality. A physics data taking to early 2019 to assess the detector performance and beam line quality. A physics data taking to the second part of 2020.

3. Constraints for $m_{A'} < 1$ MeV

Strong constraints exist for the invisible massive dark photon in the region $m_{A'} < 1$ MeV. They come from different sources:

- Atomic and nuclear experiments: These experiments aim to detect modifications of the Coulomb force (as discussed in (Bartlett et al., 1970)) due to the dark photon. Corrections in Rydberg atoms, Lamb shift and hyperfine splitting in atomic hydrogen have been translated into bounds on the massive dark photon mixing parameter (Jaekel and Roy, 2010). The results of the TEXONO neutrino experiment (Deniz et al., 2010) have been interpreted in terms of dark photon parameters in (Danilov et al., 2019);

- Axion-like particles and helioscopes: Experiments of light shining through a wall (LSW) for axions and axion-like particles can be adapted to the dark photon and limits can accordingly be estimated (Ehret et al., 2010). The same phenomenon has been used in the experiment CROWS (Betz et al., 2013) at CERN. The CAST result, on the flux of axion-like particles from the Sun (Helioscope), can be translated (Redondo, 2008) into a bound on the massive dark photon parameters. The same is true for XENON10, whose data set provides further limits (An et al., 2013a) and the results from the experiment SHIPS (Schwarz et al., 2015);

- Astrophysics: The non-observation of anomalous energy transport (by the mechanism discussed in section II.A) in stars on the horizontal branch (HB), red giants (RG) and the Sun (SUN-T and SUN-L) imposes severe constraints on the mixing parameter of the massive dark photon. Mixing effects are important in these processes for both the longitudinal (L) and transverse (T) modes and one must use thermal field theory (An et al., 2013b; Hardy and Lasenby, 2017; Redondo and Raffelt, 2013). The dark photon partakes of the plasmon modes (see appendix C) in an effective mixing with the ordinary photon proportional to its mass (and vanishing as it goes to zero);

- Cosmology: The oscillation between the ordinary and the massive dark photon $\gamma \rightarrow A'$ induces deviations on the black body spectrum (as measured by COBE/FIRAS (Fixsen et al., 1996)) in the cosmic microwave background. This effect depends on the effective plasma mass of the dark photon and it is enhanced when this mass is equal to $m_{A'}$. The bound depicted in Fig. 12 follows the most recent evaluation (Caputo et al., 2020b; Garcia et al., 2020; Witte et al., 2020)—which includes inhomogeneities in the plasma mass—for values $m_{A'} < 10^{-15}$ MeV, and (McDermott and Witte, 2020; Mirizzi et al., 2009) for larger values.

Even stronger constraints can be derived under the assumption that the dark photon is itself the dark matter. The combination (in order of increasing values of $m_{A'}$) of
FIG. 12 Current limits on massive dark photon for $m_{A'} < 1$ MeV. Top plot: Bounds from cosmology (COBE/FIRAS (Caputo et al., 2020b; Garcia et al., 2020; McDermott and Witte, 2020; Witte et al., 2020), light through a wall (LSW) (Ehret et al., 2010), CROWS (Betz et al., 2013), CAST (Redondo, 2008), XENON10 (An et al., 2013a), SHIPS (Schwarz et al., 2015), TEXONO (Danilov et al., 2019), atomic experiments (Coulomb, Rydberg and atomic spectra (Jaeckel and Roy, 2010)) and astrophysics: Solar lifetime (SUN-T and SUN-L), red giants (RG), horizontal branches (HB) (An et al., 2013b; Hardy and Lasenby, 2017; Redondo and Raffelt, 2013)). Additional limits under the assumption that the dark photon is the dark matter: The curve “Dark Matter” includes the combination of the constraints from the references discussed in the main text. Bottom plot: Zoom in the range $10^{-8}$ MeV $\lesssim m_{A'} \lesssim 1$ MeV and $10^{-17} \lesssim \varepsilon \lesssim 10^{-9}$. Results from dark matter direct detection experiments (XENON10 and XENON100 ((An et al., 2015) based on XENON10 (Angle et al., 2011) and XENON100 (Aprile et al., 2014) data), XENON1T (Aprile et al., 2019); DAMIC (Aguilar-Arevalo et al., 2017); SuperCDMS (Aralis et al., 2020); CDEX-10 (She et al., 2020); EDELWEISS-III (Armengaud et al., 2018); SENSEI (Abramoff et al., 2019); XMASS (Abe et al., 2018); FUNK (Andrianavalomahefa et al., 2020)).

- astrophysical bounds on dwarf galaxies (Wadekar and Farrar, 2019). These limits apply for values of
yield a series of limits on the upper value of $m_A$. These limits are depicted together by the curve labelled “Dark Matter” in Fig. 12.

In addition, there are limits from:

- Dark matter direct detection experiments: These experiments are part of the ongoing search for dark matter through its direct detection. Data from XENON10/XENON100 (An et al., 2015) and XENON100 (Aprile et al., 2014), XENON1T (Aprile et al., 2018), DAMIC (Aguilar-Arevalo et al., 2017), SuperCDMS (Aralis et al., 2020), CDEX-10 (She et al., 2020), EDELWEISS-III (Armengaud et al., 2018), SENSEI (Abramoff et al., 2019), XMASS (Abe et al., 2018) and FUNK (Andrianavalomahefa et al., 2020) can be used to constrain the massive dark photon parameters;

- Haloscopes: Searches with microwave cavities for relic axion converting to photons (Sikivie, 1983) can be translated into limits (not shown in Fig. 12) on the dark photon parameter $\varepsilon$ to be less than $10^{-13}$, $10^{-15}$ in the range around $10^{-11}$-10$^{-12}$ MeV (Arias et al., 2012).

The limits from dark-matter direct detection are shown in the bottom half of Fig. 12.

Some of the limits on the right side of the top plot of Fig. 12 are the continuation of the corresponding left side of the limits in Fig. 11. The two figures are back-to-back at $m_A$ = 1 MeV thus covering the full range of the dark-photon masses.

B. Constraints on $y$ and $m_\chi$

If the dark sector states into which the invisible dark photon decays are taken to be dark matter, there are new limits involving also the coupling strength $\alpha_\gamma$ and the connection to the direct-detection searches for dark matter. As discussed in section I.C, the best way to plot the experimental limits in this case is in terms of the yield variable $y$, defined in Eq. (28), and the dark matter mass $m_\chi$.

The corresponding limits strongly depend on the nature of the dark-matter state $\chi$ because the velocity dependence of the averaged cross sections. In the case of Dirac fermions, Planck data (Ade et al., 2016) rule out sub-GeV dark matter because of their too large annihilation rate at the cosmic microwave background epoch. For this reason, pseudo-Dirac fermions and scalars, which have velocity suppressed annihilation cross sections, are usually studied.

The current bounds and future perspectives in the plane $y$ versus dark matter mass are shown in Fig. 13 under the hypothesis that the dark matter is a scalar particle and for a specific choice of $\alpha_\gamma$ ($\alpha_\gamma = 0.1$) and the ratio between the mediator and the dark matter masses ($m_A/m_\chi = 3$). In these plots, the lower limit for the thermal relic density is also shown, under that hypothesis is provided by accelerator-based experiments are produced in relativistic regime and the strength of the interactions with light mediators and SM particles is only fixed by thermal freeze-out.

Current bounds come from the same experiments using missing energy/missing momentum techniques contributing to the $\{\varepsilon, m_A\}$ sensitivity plot (BaBar and NA64(e)) with the addition of the re-interpretation of data from old neutrino experiments (E137 (Batell et al., 2014) and LSND (deNiverville et al., 2011)) and results from current neutrino experiments (MiniBooNE (Aguilar-Arevalo et al., 2018)) exploiting dark matter scattering on nucleons and/or electrons. Bounds can also be derived by using a superfluid He-4 detector, as shown in (Caputo et al., 2020a), but they lie at the margin of the range included in Fig. 13.

Future initiatives that could explore a still uncovered parameter space in the plane $\{y, m_\chi\}$ for dark matter masses below 1 GeV are all those that have sensitivity in the plane $\{\varepsilon, m_A\}$ and, in addition, accelerator-based and dark matter direct detection experiments exploiting dark matter scattering against the nucleons and/or electrons. Accelerator-based experiments are SHiP at CERN (Anelli et al., 2015), and BDX at JLAB (Battaglieri et al., 2016) and SBND (Antonello et al., 2015) at FNAL as explained below.

- **BDX at JLAB:** The Beam Dump eXperiment (BDX) (Battaglieri et al., 2016) is aiming to detect light dark matter $\chi$ produced in the interaction of...
FIG. 13 Existing limits (top, filled areas) and future sensitivities for existing or proposed experiments (bottom, colored curves) for massive dark photon for $m_{A'} > 1$ MeV in the plane of the yield variable $y$ as a function of dark matter mass $m_\chi$ for an elastic scalar dark matter particle. Top plot: Limits from BaBar (Lees et al., 2017), NA64(e) (Banerjee et al., 2019), reinterpretation of the data from E137 (Batell et al., 2014) and LSND (deNiverville et al., 2011); result from MiniBooNE (Aguilar-Arevalo et al., 2018); interpretation in the dark photon framework of data from XENON10 (Angle et al., 2011), XENON100 (Aprile et al., 2016), CRESST-II (Angloher et al., 2016), and SENSEI operating with a 2 g detector in the NUMI tunnel (Barak et al., 2020). Bottom plot: Projections for SHiP (Anelli et al., 2015), BDX (Battaglieri et al., 2016), SBND (Antonello et al., 2015), LDMX@CERN (Akesson et al., 2018b; Raubenheimer et al., 2018), SENSEI with a proposed 100 g detector operating at SNOLAB (Battaglieri et al., 2017), and SuperCDMS at SNOLAB (Agnese et al., 2017). The bottom plot is revised from (Beacham et al., 2020).

an intense (100 $\mu$A) 10 GeV electron beam with a dump. The experiment is sensitive to elastic dark matter scattering $e^- \chi \rightarrow e^- \chi$ in the detector after production in $e^-Z \rightarrow e^-Z A'(A' \rightarrow \chi \chi)$.

- SBND is planned to be installed at the 8 GeV pro-
ton Booster Neutrino Beamline at FNAL about 470 m downstream of the beam dump (Antonello et al., 2015). The dark matter beam is primarily produced via pion decays out of collisions from the primary proton beam, and identified via dark-matter-nucleon or dark-matter-electron elastic scattering in a LAr-based detector. SBND is expected to improve upon MiniBooNE by more than an order of magnitude with $6 \times 10^{20}$ protons-on-target.

Also dark matter direct-detection experiments with sensitivity below 1 GeV mass contribute to this plot. These are:

- **SENSEI** is a direct detection experiment (Tiffenberg et al., 2017) that will be able to explore dark matter candidates with masses in the 1 eV and few GeV range, by detecting the signal released in dark-matter-electron scattering interactions in a fully depleted silicon CCD. A 2-gram detector is already operating in the NUMI access tunnel (Barak et al., 2020). A larger project (100 grams) can be deployed at SNOLAB if funding is obtained (Battaglieri et al., 2017).

- **CRESST-II** (Angloher et al., 2016) uses cryogenic detectors to search for nuclear recoil events induced by elastic scattering of dark-matter particles in CaWO$_4$ crystals. Because of its low-energy threshold, the sensitivity to dark matter was extended in the sub-GeV region. Current bounds are derived from a dataset corresponding to 52 kg live days.

- **XENON10/XENON100** Dark-matter-electron scattering searches have already illustrated their potential, probing down to $m_\chi > 5$ MeV (Essig et al., 2012a, 2017) using XENON10 data (Angle et al., 2011) sensitive to single electrons and down to $m_\chi > 35$ MeV using XENON100 data (Aprile et al., 2016).

**IV. CONCLUDING REMARKS**

In the past 50 years it has been assumed that physics beyond the SM interacted through (at least) some of the same gauge interactions of the SM. The minimal supersymmetric SM and weakly interacting massive dark matter are the two preeminent and most influential models based on this paradigm.

This program is now running out of some of the initial momentum because of the lack of the discovery of new particles. In the absence of new states, the many parameters, for instance, of the minimal supersymmetric SM are working against its usefulness, as a foil for the SM, in mapping possible experimental discrepancies.

In more recent times—mostly under the influence of this lack of any real signal of the breaking up of the SM—a more general scenario has been attracting increasing interest. Matter beyond the SM is part of a new sector which is dark because it does not interact through the SM gauge interactions. The dark sector may contain a wealth of physics with many particles (some of which are dark matter) and interactions.

From our side, in the visible world, we may glimpse this dark sector through a portal. If it exists, this portal can take various forms depending on the spin of the mediator. We have reviewed the vector case in which the portal arises from the kinetic mixing between the SM electric (or hyper) charge gauge group and an $U(1)$ gauge symmetry of the dark sector.

The discovery of the dark photon is by far more interesting than finding just a new particle because, if found, this new gauge boson would be the harbinger of a new interaction and of the existence of a whole new sector of elementary particles.

Past and current experiments have already restricted an important part of the space of the parameters of the vector portal, both for the massless and the massive dark photon. Compared to other searches for models beyond the SM, the parameters are fewer and the signatures more easily interpreted.

We are now on the verge of a new wave of experiments aiming at further closing the windows left still open in the interaction between ordinary matter and the dark photon.

The constraints in the massless case seem to relegate the possible detection of the dark photon to very large values of the effective scale $\Lambda$ in the dark dipole interaction, as we discuss in section II.A. Exploring physics at such a large energy scale requires the high sensitivity that can only be achieved either in future lepton colliders (where the scaling with the energy of the dark dipole operator will also enhance its contribution) or in searches for rare flavor-changing decays like those of the Kaon and $B$-mesons systems.

The constraints in the case of the massive dark photon have left open two important regions in the parameter space. The first one is for the visible dark photon with masses around 100 MeV or larger and mixing parameter between $10^{-6}$ and $10^{-4}$. Many future experiments aim at looking into this range, as we review in section III.A.1. If also this window will be closed, it means that the already feeble interaction of the vector portal is very weak indeed. The second window is for the invisible dark photon with a very light mass and a mixing parameter of order $O(10^{-8})$ or even lighter and with smaller mixing parameter, as discussed in section III.A.2 and III.A.3. These two latter regions are of great interest for astrophysics and cosmology and a very active area of speculations.

No single experiment or experimental approach is sufficient alone to cover the large parameter space in terms of
masses and couplings that dark photon models suggest: Synergy and complementarity among a great variety of experimental facilities are paramount, calling for a broad collaboration across different communities.

ACKNOWLEDGMENTS

The digital inclusion of some of the experimental limits in the figures was done by means of WebPlotDigitizer\(^4\). Numerical results for some of the dark-photon limits were provided by the DarkCAST\(^5\) package (Ilten et al., 2018).

We would like to thank Tram Acuña, Andrea Celentano, Monica D’Onofrio, Angelo Esposito, Oliver Fischer, Phil Ilten, Sam McDermott, Maxim Pospelov, Diego Redigolo, Josh Ruderman, Piero Ullio, Alfredo Urbano and Mike Williams for useful discussions and suggestions.

MF is affiliated to the Physics Department of the University of Trieste and the Scuola Internazionale Superiore di Studi Avanzati (SISSA), Trieste, Italy. MF and EG are affiliated to the Institute for Fundamental Physics of the Universe (IFPU), Trieste, Italy. The support of all these institutions is gratefully acknowledged. EG thanks the Department of Theoretical Physics of CERN for its kind hospitality during the preparation of this work.

Appendix A: Dark sector portal

In this appendix we give the Lagrangian for the four portals mentioned in the introduction as well as a minimal bibliography to provide their context within the search of phenomenological models beyond the Standard Model.

The dark sector is assumed to interact with the visible, SM sector through relevant operators of dimension four and five (and possibly sub-leading higher-order operators).

These portals are classified according to the spin of the mediator field. We can have:

- **dark photon** (spin 1): The portal operator arises from the kinetic is mixing between the SM photon field strength \(F_{\mu\nu}\) and a dark photon \(F'_{\mu\nu}\):

\[
\frac{\varepsilon}{2} F_{\mu\nu} F'_{\mu\nu},
\]

it is an operator of dimension four. It is assumed that the dark photon is the main carrier of the interaction among the dark sector states.

The existence of an independent \(U(1)\) group symmetry was originally proposed, in the context of supersymmetric theories, in (Fayet, 1980, 1981) and, more in general, in (Georgi et al., 1983; Okun, 1982);

- **axion** (spin 0): The operator comes from the interaction between a pseudo-scalar, the axion \(a\), and the SM photon and fermions \(\psi\):

\[
\frac{a}{f_a} F_{\mu\sigma} \tilde{F}^{\mu\sigma} + \frac{1}{f_a} \partial_\mu a \bar{\psi} \gamma^\mu \gamma_5 \psi,
\]

with operators of dimension five: the physics of this portal is based on that of the axion and related to the strong CP problem as well as axion dark matter. In many cases, the portal is generalized to an axion-like particle (ALP) with similar couplings but without the constraints of the QCD axion. The parameters of the portal are two: the mass \(m_a\) of the axion, or the ALP and the scale \(f_a\). Often the ALP is the only member of the dark sector of these models.

The original axion emerged from addressing (Peccei and Quinn, 1977) the strong CP problem induced by instantons in QCD. Light pseudo-scalar bosons are found in many models of physics beyond the Standard Model;

- **scalar** (spin 0): Interaction between a scalar \(S\) and the SM Higgs boson \(H\):

\[
(\mu S + \lambda S^2)H^\dagger H,
\]

in this case the operators are of dimension three and four. The experimental limits are often expressed in terms of the two parameters \(\nu\) and the mass \(m_S\) of the scalar singlet, and neglecting the quartic coupling \(\lambda\). In most models, the dark sector states have Yukawa-like interactions with the scalar \(S\).

The idea of a scalar singlet interacting with the Higgs boson originated within the framework of the next-to-minimal supersymmetric Standard Model (Derendinger and Savoy, 1984; Ellis et al., 1989; Frere et al., 1983) and developed independently in (Binoth and van der Bij, 1997; Patt and Wilczek, 2006; Silveira and Zee, 1985);

- **sterile neutrino** (spin 1/2): Interaction between a heavy fermion \(N\), which is a SM singlet, the SM Higgs boson and the SM fermions \(L\):

\[
y_N \bar{H} N,
\]

with, again, an operator of dimension four. The existence of heavy lepton-like fermions is suggested by neutrino see-saw models and the possible origin of baryon-number asymmetry in the leptonic sector. The experimental searches are framed in terms of the parameter \(y_N\) and the mass of the

\footnotesize

\(^4\) Website: https://automeris.io/WebPlotDigitizer

\(^5\) Website: https://gitlab.com/philten/darkcast
heavy fermion $N$. The sterile neutrino can be the only member of the dark sector or be one among many other dark fermions.

The structure of the neutrino portal closely follows that of the see-saw mechanism (Gell-Mann et al., 1979; Minkowski, 1977; Mohapatra and Senjanovic, 1981; Yanagida, 1980)—which was introduced to generate small masses for the neutrinos—and, more in general, left-right symmetric models (Mohapatra and Pati, 1975a,b; Pati and Salam, 1974; Senjanovic and Mohapatra, 1975).

More details on the various portals can be found in the same references cited in the introduction: (Alekhin et al., 2016; Alexander et al., 2016; Beacham et al., 2020; Curciarello, 2016; Deliyergiyev, 2016; Essig et al., 2013; Hewett et al., 2012; Raggi and Kozhinharov, 2015).

**Appendix B: Boltzmann equation and relic density**

This appendix includes a short summary of some results necessary to follow the discussion in the main text on cosmology. We follow the excellent review (Bauer and Plehn, 2019).

The rate $\Gamma$ for a the interaction between two particles is given as

$$\Gamma = n \sigma v,$$

the product of the corresponding cross section $\sigma$ times the number density of the particles partaking $n$, times the their relative velocity $v$.

This process proceeds as long as the rate is larger than the Hubble constant

$$H(T) = \frac{\pi \sqrt{g_*}(T) T^2}{\sqrt{90} m_{Pl}},$$

where $m_{Pl}$ the Planck mass and $g_*(T)$ is the number of effective degrees of freedom at the given temperature is given by

$$g_*(T) = \sum_{\text{bosons}} g_b \left( \frac{T}{T_*} \right)^4 + \frac{7}{8} \sum_{\text{fermions}} g_f \left( \frac{T_f}{T} \right)^4,$$

where $g_{b,f}$ is the number of degrees of freedom of the corresponding particle. The value of the function $g_*(T)$ goes from 106.5 above the EW phase transition to 3.38 at temperature around 0.1 MeV.

After $\Gamma < H$, the particles are decoupled and their number density frozen.

The number density at the equilibrium at a given temperature $T$ (for $k_B = 1$) is given by

$$n_{eq}(T) = g_* \int \frac{d^3 p}{(2\pi)^3} \frac{1}{e^{E/T} \pm 1} = \begin{cases} \frac{g_*}{(2\pi)^{3/2} \sqrt{g_0 T^3}} e^{-m/T} & \text{non-relativistic \quad } (T \ll m) \\ \frac{\zeta(3) \pi^2}{2} g_* T^3 & \text{relativistic bosons \quad } (T \gg m) \\ \frac{3 \zeta(3)}{4} \pi^2 g_* T^3 & \text{relativistic fermions \quad } (T \gg m) \end{cases}$$

where $\zeta(3) \approx 1.2$.

The number density $n(t)$ of a weakly interacting, massive particle $\chi$ at a certain time $t$ in the evolution of the Universe is computed by means of the Boltzmann equation

$$\dot{n}(t) + 3H(t)n(t) = -\langle \sigma_{\chi\chi\rightarrow ff} v \rangle \left( n^2(t) - n_{eq}^2(t) \right),$$

where $H(t)$ is the Hubble constant and $\langle \sigma_{\chi\chi\rightarrow ff} v \rangle$ is the thermal average of the cross section for a pair of the particles $\chi$, with relative velocity $v = (s - 4m_\chi^2)/m_\chi^2$, to annihilate into SM fermions $f$; this term depletes the density as the particles $\chi$ turns into SM fermions. The thermal average is defined as

$$\langle \sigma_{\chi\chi\rightarrow ff} v \rangle = \int_{-\infty}^{4m_\chi^2} ds s (s - 4m_\chi^2) K_1 \left( \frac{s}{2m_\chi^2} \right) \sigma_{\chi\chi\rightarrow ff}$$

for the $s$-wave and

$$\frac{16}{\pi} \int_{-\infty}^{4m_\chi^2} ds s (s - 4m_\chi^2) K_1 \left( \frac{s}{2m_\chi^2} \right) \sigma_{\chi\chi\rightarrow ff},$$

for the $p$-wave. The leading term is $s_0$ for the dark sector Dirac fermions interacting through the dark photon, in both the $s$- and $t$-channel.

Eq. (B5) is usually re-written in terms of the function $Y(t) = n(t)/T^3$ and the variable $x = m_\chi/T = \sqrt{2TH(T = m_\chi)}$ as

$$\frac{dY}{dx} = -\frac{\lambda(x)}{x^2} \left[ Y^2(x) - Y^2_{eq} \right],$$

with

$$\lambda(x) = \frac{m_\chi^3 \langle \sigma_{\chi\chi\rightarrow ff} v \rangle}{H(T = m_\chi)},$$

where $K_1$ and $K_2$ are the Bessel function of second kind. It is usually computed after expanding

$$\langle \sigma_{\chi\chi\rightarrow ff} v \rangle = \langle s_0 + s_1 v^2 + O(v^4) \rangle$$

with $s_0$ the cross section in the $s$-wave and $s_1$ the first correction in the $p$-wave. The leading term is $s_0$ for the dark sector Dirac fermions interacting through the dark photon, in both the $s$- and $t$-channel.
and in this form numerically solved.

Eq. (B8) can be solved analytically by dropping the second term \( Y_{eq}^2(x) \)—which is small because decreasing like \( e^{-x} \)—approximating

\[
\langle \sigma_{\chi \chi \rightarrow ff} v \rangle = \sigma_{\chi \chi \rightarrow ff} v + O(v^2),
\]

(B10)

where \( v = \sqrt{2/x} \) and writing

\[
\lambda(x) = \frac{180 m_p m_\chi}{\pi \sqrt{g_{*d}}} \sigma_{\chi \chi \rightarrow ff}
\]

(B11)

by means of

\[
H(T = m_\chi) = \frac{\pi \sqrt{g_{*d} m_\chi}}{90} m_\chi / m_{Pl}.
\]

(B12)

The solution for \( x' \) larger than decoupling temperature \( x_d \) is

\[
Y(x') = \frac{x_d}{\lambda}.
\]

(B13)

This quantity is related to the relic density

\[
\rho_\chi = m_\chi n(x') = m_\chi^4 \frac{Y(x')}{28 x_d}
\]

(B14)

or, in terms of the normalized quantity \( \Omega_\chi = \rho_\chi / \rho_c \) as

\[
\Omega_\chi h^2 \approx 0.12 \frac{x_d \sqrt{m_\chi}}{23} \frac{1.7 \times 10^{-9} \text{GeV}^{-2}}{\langle \sigma_{\chi \chi \rightarrow ff} v \rangle}
\]

\[
\approx 2.5 \times 10^{-10} \text{GeV}^{-2}
\]

(B15)

which provides the relationship between relic density and the annihilation cross section.

Appendix C: Thermal field theory

The energy loss rate \( Q \) (energy per volume and unit time) for the emission of a pseudoscalar particle (the axion) in a process with matrix element, which is computed in the vacuum, is given by

\[
Q = \prod_{i=1} d^3p_i f_i(E_i) \prod_{f=1} d^3p_f f_f(E_f) \int \frac{d^4p_a}{2\omega_a(2\pi)^3} \langle 1 \pm f_f(E_f) \rangle \frac{1}{3} \sum_{\text{spin and pol.}} |\mathcal{M}|^2(2\pi)^4 \delta^4 \left( \sum p_i - \sum p_f - p_a \right),
\]

(C1)

in section II.A, the corresponding luminosity can be computed as

\[
L = \int dV Q e^{-\tau},
\]

(C3)

where \( \tau \) is an attenuation factor taking into account the optical depth of the emission, and compared to the observational data.

When the emitted particle mixes with the ordinary photon, the approach above of computing the matrix element in the vacuum is no longer a reliable approximation and the full thermal field theory must be used. We follow (Braaten and Segel, 1993) in giving the essential formulas.

The electromagnetic polarization tensor is given by

\[
\Pi^{\mu\nu}(k) = 16\pi a \int \frac{d^3p}{(2\pi)^3} \frac{1}{2E} \left[ n_e(E) + n_\nu(E) \right] \frac{p \cdot k (p^\mu k^\nu + k^\mu p^\nu) - k^2 p^\mu p^\nu - (p \cdot k)^2 g^{\mu\nu}}{(p \cdot k)^2 - (k^2)^2/4},
\]

(C4)

longitudinal polarizations are defined as

\[
\Pi_T(\omega, k) = \frac{1}{2} (\delta^{ij} - k^i k^j) \Pi^{ij}(\omega, k)
\]

(C5)
and
\[ \Pi_L(\omega, k) = \Pi^0(\omega, k). \]  
(C6)

The effective propagator of the photon (in the Coulomb gauge) has components
\[ D^{00}(\omega, k) = \frac{1}{k^2 - \Pi_L(\omega, k)} \]  
(C7)

and
\[ D^{ij}(\omega, k) = \frac{1}{k^2 - \Pi_T(\omega, k)} (\delta^{ij} - k^i k^j). \]  
(C8)

The dispersion relationships are defined by the solutions of the equations
\[ \omega^2 = k^2 + \Pi_T(\omega_T, k) \quad \text{and} \quad \omega^2 = \frac{\omega^2}{k^2} + \Pi_T(\omega_L, k) \]  
(C9)

In the degenerate limit, the distribution functions in Eq. (C2) reduce to step functions at the Fermi momentum \( p_F = \sqrt{3\pi^2 n_e} \) and we have
\[ \Pi_T(\omega, k) = \frac{3\omega^2}{2v_F^2 k^2} \left( 1 - \frac{\omega^2 - v_F^2 k^2}{2v_F^2 k^2} \log \frac{\omega + v_F k}{\omega - v_F k} \right) \]  
(C10)

and
\[ \Pi_L(\omega, k) = \frac{3\omega}{2v_F^2 k^2} \left( \frac{\omega}{2v_F k} \log \frac{\omega + v_F k}{\omega - v_F k} - 1 \right) \]  
(C11)

where \( \omega_F = 4\alpha P_F^2 v_F/3\pi \) is the plasma frequency.

The energy loss rate \( Q \) (energy per volume and unit time) is written in terms of the imaginary part of the polarization of the photon in the medium of charged particles. The contribution of the longitudinal and transverse modes is obtained (by the optical theorem) as
\[ Q = -\int \frac{d^3k}{(2\pi)^3} \frac{\text{Im} \Pi_L(\omega, k) + \text{Im} \Pi_T(\omega, k)}{\omega(\epsilon^2/T - 1)}, \]  
(C12)

from which the luminosity in Eq. (C3) can be computed and compared with the astrophysical limit of interest. The expression for \( \text{Im} \Pi_L \) for the massive dark photon can be found in (An et al., 2013b; Redondo and Raffelt, 2013) and (Hardy and Lasenby, 2017) where plasma effects are also included.

REFERENCES

Morad Aaboud, et al. (ATLAS) (2016), “Search for new phenomena in events with a photon and missing transverse momentum in pp collisions at \( \sqrt{s} = 13 \) TeV with the ATLAS detector,” JHEP 06, 059, arXiv:1604.01306 [hep-ex].

Morad Aaboud, et al. (ATLAS) (2018), “Search for squarks and gluinos in final states with jets and missing transverse momentum using 36 fb\(^{-1}\) of \( \sqrt{s} = 13 \) TeV pp collision data with the ATLAS detector,” Phys. Rev. D97 (11), 112001, arXiv:1712.02332 [hep-ex].

Georges Aad, et al. (ATLAS) (2014), “Search for long-lived neutral particles decaying into lepton jets in proton-proton collisions at \( \sqrt{s} = 8 \) TeV with the ATLAS detector,” JHEP 11, 088, arXiv:1409.0746 [hep-ex].

Georges Aad, et al. (ATLAS) (2016), “A search for prompt lepton-jets in pp collisions at \( \sqrt{s} = 8 \) TeV with the ATLAS detector,” JHEP 02, 062, arXiv:1511.05542 [hep-ex].

Roel Aaij, et al. (LHCb) (2018a), “Physics case for an LHCb Upgrade II - Opportunities in flavour physics, and beyond, in the HL-LHC era,” arXiv:1808.08865 [hep-ex].

Roel Aaij, et al. (LHCb) (2018b), “Search for Dark Photons Produced in 13 TeV pp Collisions,” Phys. Rev. Lett. 120 (6), 061801, arXiv:1710.02867 [hep-ex].

Roel Aaij, et al. (LHCb) (2020), “Search for \( A' \to \mu^+ \mu^- \) Decays,” Phys. Rev. Lett. 124 (4), 041801, arXiv:1910.06920 [hep-ex].

Laura G. van den Aarsen, Torsten Bringmann, and Christoph Pfommer (2012), “Is dark matter with long-range interactions a solution to all small-scale problems of A CDM cosmology?” Phys. Rev. Lett. 109, 231301, arXiv:1205.5809 [astro-ph.CO].

G. Abbiendi, et al. (OPAL) (1999), “Search for anomalous photonic events with missing energy in \( e^+e^- \) collisions at \( \sqrt{s} = 130-\text{GeV}, 136-\text{GeV}, \) and 183-\text{GeV},” Eur. Phys. J. C8, 23-40, arXiv:hep-ex/9810021 [hep-ex].

L. F. Abbott, and P. Sikivie (1983), “A Cosmological Bound on the Invisible Axion,” Phys. Lett. 120B, 133-136.

K. Abe, et al. (XMASS) (2018), “Search for dark matter in the form of hidden photons and axion-like particles in the XMASS detector,” Phys. Lett. B787, 153-158, arXiv:1807.08516 [astro-ph.CO].

S. A. Abel, and B. W. Schofield (2004), “Brane anti-brane kinetic mixing, millicharged particles and SUSY breaking,” Nucl. Phys. B685, 150-170, arXiv:hep-th/0311051 [hep-th].

Steven A. Abel, Joerg Jaeckel, Valentin V. Khoze, and Andreas Ringwald (2008), “Illuminating the Hidden Sector of String Theory by Shining Light through a Magnetic Field,” Phys. Lett. B666, 66-70, arXiv:hep-ph/0608248 [hep-ph].

Orr Abramoff, et al. (SENSEI) (2019), “SENSEI: Direct-Detection Constraints on Sub-GeV Dark Matter from a Shallow Underground Run Using a Prototype Skipper-CCD,” Phys. Rev. Lett. 122 (16), 161801, arXiv:1901.10478 [hep-ex].

M. Acciarri, et al. (L3) (1999), “Single and multiphoton events with missing energy in \( e^+e^- \) collisions at \( \sqrt{s} = 189-\text{GeV},” Phys. Lett. B470, 268-280, arXiv:hep-ex/9910009 [hep-ex].

Lotty Ackerman, Matthew R. Buckley, Sean M. Carroll, and Marc Kamionkowski (2009), “Dark Matter and Dark Radiation,” Phys. Rev. D79, 023519, [277(2008)], arXiv:0810.5126 [hep-ph].

Jan Tristram Acuña, Marco Fabbrichesi, and Piero Ullio (2020), “Phenomenological consequences of an interacting multicomponent dark sector,” arXiv:2005.04146 [hep-ph].

P.A.R. Ade, et al. (Planck) (2016), “Planck 2015 results. XIII. Cosmological parameters,” Astron. Astrophys. 594, A13, arXiv:1502.01589 [astro-ph.CO].

S. Adler, et al. (E787) (2002), “Further evidence for the decay \( K^+ \to \pi^+ \nu \bar{\nu} \),” Phys. Rev. Lett. 88, 041803, arXiv:hep-ex/0111091 [hep-ex].

P.H. Adrian, et al. (IPS) (2018), “Search for a dark photon in electroproduced \( e^+e^- \) pairs with the Heavy Photon
Search experiment at JLab,” Phys. Rev. D 98 (9), 091101, arXiv:1807.11530 [hep-ex].
R. Agnese, et al. (SuperCDMS) (2017), “Projected Sensitivity of the SuperCDMS SNOLAB experiment,” Phys. Rev. D 95 (8), 082002, arXiv:1610.00006 [physics.ins-det].
Prateek Agrawal, Francis-Yan Cyr-Racine, Lisa Randall, and Jakub Scholtz (2017), “Make Dark Matter Charged Again,” JCAP 1705 (05), 022, arXiv:1610.04611 [hep-ph].
F. del Aguila, M. Masip, and M. Perez-Victoria (1995), “Physical parameters and renormalization of U(1)-a x U(1)-b models,” Nucl. Phys. B456, 531–549, arXiv:hep-ph/9507455 [hep-ph].
A. Aguilar-Arevalo, et al. (DAMIC) (2017), “First Direct-Detection Constraints on eV-Scale Hidden-Photon Dark Matter with DAMIC at SNOLAB,” Phys. Rev. Lett. 118 (14), 141803, arXiv:1611.03066 [astro-ph.CO].
S. Alekhin, et al. (MiniBooNE) (2018), “Dark Matter Search in Nucleon, Pion, and Electron Channels from a Proton Beam Dump with MiniBooNE,” Phys. Rev. D 98 (11), 112004, arXiv:1807.06137 [hep-ex].
J.K. Ahn, et al. (KOTO) (2019), “Search for the $K_L \rightarrow \pi^0\nu\bar{\nu}$ and $K_L \rightarrow \pi^0\pi^0$ decays at the J-PARC KOTO experiment,” Phys. Rev. Lett. 122 (02), 021802, arXiv:1810.09655 [hep-ex].
Torsten Akesson, et al. (2018a), “Dark Sector Physics with a Primary Electron Beam Facility at CERN,” .
Torsten Akesson, et al. (LDMX) (2018b), “Light Dark Matter eXperiment (LDMX),” arXiv:1808.05219 [hep-ex].
Sergey Alekhin, et al. (2016), “A facility to Search for Hidden Particles at the CERN SPS: the SHiP physics case,” Rept. Prog. Phys. 79 (12), 124201, arXiv:1504.04855 [hep-ph].
Jim Alexander, et al. (2016), “Dark Sectors 2016 Workshop: Community Report,” arXiv:1608.08632 [hep-ph].
Stephon Alexander, Evan McDonough, Robert Sims, and Nicolas Yunes (2018), “Hidden-Sector Modifications to Gravitational Waves From Binary Inspiral,” Class. Quant. Grav. 35 (23), 235012, arXiv:1808.05286 [gr-qc].
W. Allmannshofer, et al. (Belle-II) (2019), “The Belle II Physics Book,” PTEP 2019 (12), 123C01, [Erratum: PTEP2020, no.2,022901 (2020)], arXiv:1808.10567 [hep-ex].
F. Ambrosino, et al. (KLOE Project) (2019), “KLEVER: An experiment to measure $BR(K_L \rightarrow \pi^0\pi^0)$ at the CERN SPS,” arXiv:1901.03099 [hep-ex].
Haipeng An, Maxim Pospelov, and Josef Pradler (2013a), “Dark Matter Detectors as Dark Photon Helioscopes,” Phys. Rev. Lett. 111, 041302, arXiv:1304.3461 [hep-ph].
Haipeng An, Maxim Pospelov, and Josef Pradler (2013b), “New stellar constraints on dark photons,” Phys. Lett. B725, 190–195, arXiv:1302.3884 [hep-ph].
Haipeng An, Maxim Pospelov, Josef Pradler, and Adam Ritz (2015), “Direct Detection Constraints on Dark Photon Dark Matter,” Phys. Lett. B 747, 331–338, arXiv:1412.8378 [hep-ph].
A. Anastasi, et al. (KLOE-2) (2016), “Limit on the production of a new vector boson in $e^+e^- \rightarrow U\gamma, U \rightarrow \pi^+\pi^-$ with the KLOE experiment,” Phys. Lett. B757, 356–361, arXiv:1603.06086 [hep-ex].
A. Andrianovamalahaefa, et al. (FUNK Experiment) (2020), “Limits from the Funk Experiment on the Mixing Strength of Hidden-Photon Dark Matter in the Visible and Near-Ultraviolet Wavelength Range,” arXiv:2003.13144 [astro-ph.CO].
M. Anelli, et al. (SHiP) (2015), “A facility to Search for Hidden Particles (SHiP) at the CERN SPS,” arXiv:1504.04956 [physics.ins-det].
J. Angle, et al. (XENON10) (2011), “A search for light dark matter in XENON10 data,” Phys. Rev. Lett. 107, 051301, [Erratum: Phys.Rev.Lett. 110, 249901 (2013)], arXiv:1104.3088 [astro-ph.CO].
G. Angloher, et al. (CRESST) (2016), “Results on light dark matter particles with a low-threshold CRESST-II detector,” Eur. Phys. J. C 76 (1), 25, arXiv:1509.01515 [astro-ph.CO].
M. Antonello, et al. (MicroBooNE, LAr1-ND, ICARUS-WA104) (2015), “A Proposal for a Three Detector Short-Baseline Neutrino Oscillation Program in the Fermilab Booster Neutrino Beam,” arXiv:1503.01520 [physics.ins-det].
Thomas Appelquist, Bogdan A. Dobrescu, and Adam R. Hopper (2003), “Nonexotic Neutral Gauge Bosons,” Phys. Rev. D 68, 035012, arXiv:hep-ph/0212073 [hep-ph].
E. Aprile, et al. (XENON100) (2014), “First Axion Results from the XENON100 Experiment,” Phys. Rev. D 90 (6), 062009, [Erratum: Phys. Rev.D95,no.2,029904(2017)], arXiv:1404.1455 [astro-ph.CO].
E. Aprile, et al. (XENON) (2016), “Low-mass dark matter search using ionization signals in XENON100,” Phys. Rev. D 94 (9), 092001, [Erratum: Phys.Rev.D 95, 059901 (2017)], arXiv:1605.06262 [astro-ph.CO].
E. Aprile, et al. (XENON) (2019), “Light Dark Matter Search with Ionization Signals in XENON1T,” Phys. Rev. Lett. 123 (25), 251801, arXiv:1907.11485 [hep-ex].
T. Aralis, et al. (SuperCDMS) (2020), “Constraints on Dark Photons and Axion-Like Particles from SuperCDMS Soudan,” Phys. Rev. D101 (5), 052008, arXiv:1911.11905 [hep-ex].
F. Archilli, et al. (KLOE-2) (2012), “Search for a vector gauge boson in $\phi$ meson decays with the KLOE detector,” Phys. Lett. B706, 251–255, arXiv:1110.0411 [hep-ex].
Paola Arias, Davide Cadamuro, Mark Goodsell, Jörg Jaeckel, Javier Redondo, and Andreas Ringwald (2012), “WISPy Cold Dark Matter,” JCAP 1206, 013, arXiv:1201.5902 [hep-ph].
Akiyaka Ariga, et al. (FAFER) (2019), “FAFER physics reach for long-lived particles,” Phys. Rev. D 99 (9), 095011, arXiv:1811.12522 [hep-ph].
Nima Arkani-Hamed, Douglas P. Finkbeiner, Tracy R. Slatyer, and Neal Weiner (2009), “A Theory of Dark Matter,” Phys. Rev. D 79, 051304, arXiv:0810.0713 [hep-ph].
Nima Arkani-Hamed, and Neal Weiner (2008), “LHC Signals for a SuperUnified Theory of Dark Matter,” JHEP 12, 104, arXiv:0810.0714 [hep-ph].
E. Armengaud, et al. (EDELWEISS) (2018), “Searches for electron interactions induced by new physics in the EDELWEISS-III Germanium bolometers,” Phys. Rev. D 99 (8), 082004, arXiv:1808.02340 [hep-ex].
A. V. Artamonov, et al. (BNL-E949) (2009), “Study of the decay $K^- \rightarrow \pi^+\nu\bar{\nu}$ in the momentum region $140 < P_\pi < 199$ MeV/c,” Phys. Rev. D 79, 092004, arXiv:0903.0030 [hep-ex].
K. S. Babu, Christopher F. Kolda, and John March-Russell (1998), “Implications of generalized Z - Z-prime mixing,” Phys. Rev. D57, 6788–6792, arXiv:hep-ph/9710441 [hep-ph].
D. Babusci, et al. (KLOE-2) (2013), “Limit on the production of a light vector gauge boson in phi meson decays with the KLOE detector,” Phys. Lett. B720, 111–115, arXiv:1210.3927 [hep-ex].
Bremstrahlung in Beam-Dump Data," *Phys. Lett.* **B371**, 320–326, arXiv:1311.3870 [hep-ph].

Kimberly K. Boddy, Jonathan L. Feng, Manoj Kaplinghat, and Tim M. P. Tait (2014), “Self-Interacting Dark Matter from a Non-Abelian Hidden Sector,” *Phys. Rev.* **D89** (11), 115017, arXiv:1402.3629 [hep-ph].

C. Boehm, T.A. Ensslin, and J. Silk (2004), “Can Annihilating dark matter be lighter than a few GeV?” *J. Phys.* **G 30**, 279–286, arXiv:astro-ph/0208458.

C. Boehm, and Pierre Fayet (2004), “Scalar dark matter candidates,” *Nucl. Phys. B* **683**, 219–263, arXiv:hep-ph/0305261.

Eric Braaten, and Daniel Segel (1993), “Neutrino energy loss from the plasma process at all temperatures and densities,” *Phys. Rev. D* **48**, 1478–1491, arXiv:hep-ph/9302213.

Ralf Peter Brinkmann, and Michael S. Turner (1988), “Numerical Rates for Nucleon-Nucleon Axion Bremstrahlung,” *Phys. Rev. D* **38**, 2338.

Joachim Brod, Aaron Gootjes-Dreesbach, Michele Tammaro, and Jure Zupan (2018), “Effective Field Theory for Dark Matter Direct Detection up to Dimension Seven,” *JHEP* **10**, 065, arXiv:1710.10218 [hep-ph].

A. Bross, M. Crisler, Stephen H. Pordes, J. Volk, S. Errede, and J. Wrbanek (1991), “A Search for Shortlived Particles Produced in an Electron Beam Dumps,” *Phys. Rev. Lett.* **67**, 2942–2945.

Matthew R. Buckley, and Patrick J. Fox (2010), “Dark Matter Self-Interactions and Light Force Carriers,” *Phys. Rev. D* **81**, 083522, arXiv:0911.3805 [hep-ph].

Manuel A. Buen-Abad, Gustavo Marques-Tavares, and Martin Schmaltz (2015), “Non-Abelian dark matter and dark radiation,” *Phys. Rev. D92* (2), 023531, arXiv:1505.03542 [hep-ph].

A. Caldwell, et al. (2018), “Particle physics applications of the AWAKE acceleration scheme,” arXiv:1812.11164 [physics.acc-ph].

Andrea Caputo, Angelo Esposito, Emma Geoffray, Antonio D. Polosa, and Sicun Sun (2020a), “Dark Matter Self-Interactions and Light Force Carriers,” *Phys. Rev. D* **105**, 115019, arXiv:1203.2947 [hep-ph].

Andrea Caputo, Hongwan Liu, Siddharth Mishra-Sharma, and Joshua T. Ruderman (2020b), “Dark Photon Oscillations in Our Inhomogeneous Universe,” arXiv:2002.05165 [astro-ph.CO].

Pierluca Carenza, Tobias Fischer, Maurizio Giannotti, Gang Guo, Gabriel Martinez-Pinedo, and Alessandro Mirizzi (2019), “Improved axion emissivity from a supernova via nucleon-nucleon bremsstrahlung,” *JCAP* **1910** (10), 016, arXiv:1906.11844 [hep-ph].

Eric D. Carlson (1987), “LIMITS ON A NEW U(1) COUPLING,” *Nucl. Phys. B* **286**, 378–398.

Eric D. Carlson, Marie E. Machacek, and Lawrence J. Hall (1992), “Self-interacting dark matter,” *Astrophys. J.* **398**, 43–52.

Yuk Fung Chan, Matthew Low, David E. Morrissey, and Andrew P. Spray (2012), “LHC Signatures of a Minimal Supersymmetric Hidden Valley,” *JHEP* **05**, 155, arXiv:1112.2705 [hep-ph].

Jae Hyeok Chang, Rouven Eising, and Samuel D. McDermott (2018), “Supernova 1987A Constraints on Sub-GeV Dark Sectors, Millicharged Particles, the QCD Axion, and an Axion-like Particle,” *JHEP* **09**, 051, arXiv:1803.00993 [hep-ph].

So Chigusa, Takeo Moroi, and Kazunori Nakayama (2020), “Detecting Light Boson Dark Matter through Conversion into Magnon,” arXiv:2001.10666 [hep-ph].

Xiaoayong Chu, and Basudeb Dasgupta (2014), “Dark Radiation Alleviates Problems with Dark Matter Halos,” *Phys. Rev. Lett.* **113** (16), 161301, arXiv:1404.6127 [hep-ph].

Xiaoayong Chu, Jui-Lin Kuo, and Josef Pradler (2020), “Dark sector-photon interactions in proton-beam experiments,” *Phys. Rev. D* **101**, 075035, arXiv:2001.06042 [hep-ph].

James M. Cline, Zuowei Liu, Guy Moore, and Wei Xue (2014), “Composite strongly interacting dark matter,” *Phys. Rev. D* **90** (1), 015023, arXiv:1312.3325 [hep-ph].

James M. Cline, Zuowei Liu, and Wei Xue (2012), “Millicharged Atomic Dark Matter,” *Phys. Rev. D* **85**, 101302, arXiv:1201.4858 [hep-ph].

Douglas Clowe, Marusia Bradac, Anthony H. Gonzalez, Maxim Markevitch, Scott W. Randall, Christine Jones, and Dennis Zaritsky (2006), “A direct empirical proof of the existence of dark matter,” *Astrophys. J.* **648**, L109–L113, arXiv:astro-ph/0608407 [astro-ph].

Eduardo CortinaGil, et al. (NA62) (2017), “The Beam and detector of the NA62 experiment at CERN,” JINST **12** (05), P05025, arXiv:1703.08501 [physics.ins-det].

Eduardo Cortina Gil, et al. (2019a), “ADDENDUM I TO P326 Continuation of the physics programme of the NA62 experiment,” arXiv:1901.10666 [hep-ph].

Eduardo Cortina Gil, et al. (NA62) (2019b), “First search for $K^+ \rightarrow \pi^+ \nu \bar{\nu}$ using the decay-in-flight technique,” *Phys. Rev. Lett.* **123**, 15–16, arXiv:1811.08508 [hep-ex].

Eduardo Cortina Gil, et al. (NA62) (2019c), “Search for production of an invisible dark photon in $\pi^0$ decays,” *JHEP* **05**, 182, arXiv:1903.08767 [hep-ex].

Djuna Croon, Ann E. Nelson, Chen Sun, Devin G. E. Walker, and Zhong-Xiu Xianyu (2018), “Hidden-Sector Spectroscopy with Gravitational Waves from Binary Neutron Stars,” *Astrophys. J.* **858** (1), L2, arXiv:1711.02096 [hep-ph].

Francesca Cuciricciello (2016), “Review on Dark Photon,” *Proceedings, Workshop on Flavour changing and conserving processes 2015 (FCCP2015): Anacapri, Capri Island, Italy, September 10–12, 2015*, EPJ Web Conf., **118**, 01008.

David Curtin, Rouven Eising, Stefania Gori, and Jessie Shelton (2015), “Illuminating Dark Photons with High-Energy Colliders,” *JHEP* **02**, 157, arXiv:1412.0018 [hep-ph].

Francis-Yan Cyr-Racine, and Kris Sigurdson (2013), “Cosmology of atomic dark matter,” *Phys. Rev. D* **87** (10), 103515, arXiv:1209.5752 [astro-ph.CO].

Mikhail Danilov, Sergey Demidov, and Dmitry Gorbunov (2019), “Constraints on hidden photons produced in nuclear reactors,” *Phys. Rev. Lett.* **122** (4), 041801, arXiv:1804.10777 [hep-ph].

Sacha Davidson, Steen Hannestad, and Georg Raffelt (2000), “Updated bounds on millicharged particles,” *JHEP* **05**, 003, arXiv:hep-ph/0001179 [hep-ph].

Hoeman Davoudiasl, Hye-Sung Lee, and William J. Marciano (2012), “Dark’ Z implications for Parity Violation, Rare Meson Decays, and Higgs Physics,” *Phys. Rev. D* **85**, 115019, arXiv:1203.2947 [hep-ph].

Eugenio Del Nobile, Chris Kouvaris, Paolo Panci, Francesco Sannino, and Juszi Virkajarvi (2012), “Light Magnetic Dark Matter in Direct Detection Searches,” *JCAP* **08**, 010, arXiv:1203.6652 [hep-ph].

M. A. Deliyergiyev (2016), “Recent Progress in Search for Dark Sector Signatures,” *Open Phys.* **14** (1), 281–303,
Klaus Ehret, et al.

Richard Keith Ellis, et al.

Luca Doria, Patrick Achenbach, Mirco Christmann, Achim Jurgen Engelfried (NA62) (2019), “Search for...

J. Engel, D. Seckel, and A. C. Hayes (1990), “Emission...

Bogdan A. Dobrescu (2005), “Massless gauge bosons other...

James B. Dent, Francesc Ferrer, and Lawrence M. Krauss (2012), “Constraints on Light Hidden Sector Gauge Bosons from Supernova Cooling,” arXiv:1201.2683 [astro-ph.CO].

J. Dent, J. Ferrer, and L. Krauss (2013), “Dark Matter and Lightness,” in Proceedings, 2013 Community Summer Study on the Future of U.S. Particle Physics: Snowmass on the Mississippi (CSS2013): Minneapolis, MN, USA, July 29-August 6, 2013, arXiv:1311.0029 [hep-ph].

Rouven Essig, E. Gabrielli, and B. Mele (2017), “Hunting down massless dark photons in kasonic physics,” Phys. Rev. Lett. 119 (3), 031801, arXiv:1705.03470 [hep-ph].

Rouven Essig, E. Gabrielli, and B. Mele (2018), “Z Boson Decay into Light and Darkness,” Phys. Rev. Lett. 120 (17), 171803, arXiv:1712.05412 [hep-ph].

Marco Fabbrichesi, and Emidio Gabrielli (2019), “Dark-sector physics in the search for the rare decays K^+ \rightarrow \pi^+\nu\bar{\nu} and K_L \rightarrow \pi^+\nu\bar{\nu},” arXiv:1911.03755 [hep-ph].

Marco Fabbrichesi, and Alfredo Urbano (2019), “Charged neutron stars and observational tests of a dark force weaker than gravity,” arXiv:1902.07914 [hep-ph].

Pierre Fayet (1980), “Effects of the Spin 1 Partner of the Goldstino (Gravitino) on Neutral Current Phenomenology,” Phys. Lett. B 95, 285–289.

Pierre Fayet (1981), “On the Search for a New Spin 1 Boson,” Nucl. Phys. B 187, 184–204.

Pierre Fayet (1990), “Extra U(1)'s and New Forces,” Nucl. Phys. B 347, 743–768.

Pierre Fayet (2017), “The light U boson as the mediator of a new force, coupled to a combination of Q, B, L and dark matter,” Eur. Phys. J. C77 (1), 53, arXiv:1611.05357 [hep-ph].

Daniel Feldman, Zuowei Liu, and Pran Nath (2007), “The Stueckelberg Z-prime Extension with Kinetic Mixing and Milli-Charged Dark Matter From the Hidden Sector,” Phys. Rev. D75, 115001, arXiv:hep-ph/0702123 [HEP-PH].

Jonathan L. Feng, Bartosz Fornal, Itahal Galon, Susan Gardner, Jordan Smolinsky, Tim M. P. Tait, and Philip Tanedo (2016), “Protophobic Fifth-Force Interpretation of the Observed Anomaly in the Be Nuclear Transitions,” Phys. Rev. Lett. 117 (7), 071803, arXiv:1604.07411 [hep-ph].

Jonathan L. Feng, Bartosz Fornal, Itahal Galon, Susan Gardner, Jordan Smolinsky, Tim M. P. Tait, and Philip Tanedo (2017), “Particle physics models for the 17 MeV anomaly in beryllium nuclear decays,” Phys. Rev. D 95 (3), 035017, arXiv:1608.03591 [hep-ph].

Jonathan L. Feng, Itahal Galon, Felix Kling, and Sebastian Trojanowski (2018), “ForwArd Search ExpeRiment at the LHC,” Phys. Rev. D97 (3), 035001, arXiv:1708.09389 [hep-ph].

Jonathan L. Feng, Manoj Kaplinghat, Huitzu Tu, and Hai-Bo Yu (2009), “Hidden Charged Dark Matter,” JCAP 0907,
Jonathan L. Feng, Mano) Kaplinghat, and Hai-Bo Yu (2010), “Halo Shape and Relic Density Exclusions of Sommerfeld-Enhanced Dark Matter Explanations of Cosmic Ray Elevations,” Phys. Rev. Lett. 104, 151301, arXiv:0911.0422 [hep-ph].

Jonathan L. Feng, Huitzu Tu, and Hai-Bo Yu (2008), “Thermal Relics in Hidden Sectors,” JCAP 0810, 043, arXiv:0808.2318 [hep-ph].

E. Fermi (1934), “An attempt of a theory of beta radiation. I.” Zeitschrift für Physik 88, 161.

Filip Ficek, Derek F. Jackson Kimball, Mikhail Kozlov, Nathan Leeder, Szymon Pustelny, and Dmitry Budker (2017), “Constraints on exotic spin-dependent interactions between electrons from helium fine-structure spectroscopy,” Phys. Rev. A95 (3), 032505, arXiv:1608.05779 [physics.atom-ph].

Brian D. Fields, Keith A. Olive, Tsung-Han Yeh, and Charles Young (2020), “Big-Bang Nucleosynthesis After Planck,” JCAP 2003 (03), 010, arXiv:1912.01132 [astro-ph.CO].

A. Liam Fitzpatrick, Patrick Haxton, Emanuel Katz, Nicholas Lubbers, and Yiming Xu (2013), “The Effective Field Theory of Dark Matter Direct Detection,” JCAP 1302, 004, arXiv:1203.3542 [hep-ph].

D.J. Fidcen, E.S. Cheng, J.M. Gales, John C. Mather, R.A. Shafer, and E.L. Wright (1996), “The Cosmic Microwave Background spectrum from the full COBE FIRAS data set,” Astrophys. J. 473, 576, arXiv:astro-ph/9605054.

Robert Foot (2004), “Mirror matter-type dark matter,” Int. J. Mod. Phys. D13, 2161–2192, arXiv:astro-ph/0407623 [astro-ph].

N. Fornengo, P. Panci, and M. Regis (2011), “Long-Range forces in Direct Dark Matter Searches,” Phys. Rev. D 84, 115002, arXiv:1108.4661 [hep-ph].

Anthony Fradette, Maxim Pospelov, Josef Pradler, and Adam Ritz (2014), “Cosmological Constraints on Very Dark Photons,” Phys. Rev. D90 (3), 035022, arXiv:1407.0993 [hep-ph].

J.M. Friere, D.R.T. Jones, and S. Raby (1983), “Fermion Masses and Induction of the Weak Scale by Supergavity,” Nucl. Phys. B 222, 11–19.

Emidio Gabrielli, Matti Heikinheimo, Barbara Mele, and Martti Raidal (2014), “Dark photons and resonant monophoton signatures in Higgs boson decays at the LHC,” Phys. Rev. D90 (5), 055032, arXiv:1405.5196 [hep-ph].

Emidio Gabrielli, Luca Marzola, and Martti Raidal (2017), “Radiative Yukawa Couplings in the Simplest Left-Right Symmetric Model,” Phys. Rev. D95 (3), 035005, arXiv:1611.0009 [hep-ph].

Emidio Gabrielli, Barbara Mele, Martti Raidal, and Elena Venturini (2016), “FCNC decays of standard model fermions into a dark photon,” Phys. Rev. D94 (11), 115013, arXiv:1607.05928 [hep-ph].

Emidio Gabrielli, and Martti Raidal (2014), “Exponentially spread dynamical Yukawa couplings from nonperturbative chiral symmetry breaking in the dark sector,” Phys. Rev. D89 (1), 015008, arXiv:1310.1000 [hep-ph].

Peter Galison, and Anesh Manohar (1984), “TWO Z’s OR NOT TWO Z’s?” Phys. Lett. 136B, 279–283.

Andres Arambaru Garcia, Kyrillo Bondarenko, Sylvia Ploegkinger, Josef Pradler, and Anastasia Sokolenko (2020), “Effective photon mass and (dark) photon conversion in the inhomogeneous Universe,” arXiv:2003.10465 [astro-ph.CO].

Murray Gell-Mann, Pierre Ramond, and Richard Slansky (1979), “Complex Spinors and Unified Theories,” Conf. Proc. C709027, 315–321, arXiv:1306.4669 [hep-th].

Howard Georgi, Paul H. Ginsparg, and S. L. Glashow (1983), “Photon Oscillations and the Cosmic Background Radiation,” Nature 306, 765–766.

Tony Gherghetta, Jörn Kersten, Keith olive, and Maxim Pospelov (2019), “Evaluating the price of tiny kinetic mixing,” Phys. Rev. D 100 (9), 095001, arXiv:1909.00696 [hep-ph].

Maurizio Giannotti, Igor Iarostorza, Javier Redondo, and Andreas Ringwald (2016), “Cool WISPs for stellar cooling excesses,” JCAP 1605 (05), 057, arXiv:1512.08108 [astro-ph.HE].

G.F. Giudice, P. Paradisi, and M. Passera (2012), “Testing new physics with the electron g-2,” JHEP 11, 113, arXiv:1208.6853 [hep-ph].

S. N. Gninenko (2015), “Search for invisible decays of $\pi^0, \eta, \eta'$, $K_S$ and $K_L$: A probe of new physics and tests using the Bell-Steinberger relation,” Phys. Rev. D91 (1), 015004, arXiv:1409.2288 [hep-ph].

S. N. Gninenko, and N. V. Krasnikov (2015), “Invisible $K_L$ decays as a probe of new physics,” Phys. Rev. D92 (3), 034009, arXiv:1503.01505 [hep-ph].

S.N. Gninenko (2012), “Constraints on sub-GeV hidden sector gauge bosons from a search for heavy neutrino decays,” Phys. Lett. B713, 244–248, arXiv:1204.3583 [hep-ph].

Haim Goldberg, and Lawrence J. Hall (1986), “A New Candidate for Dark Matter,” Phys. Lett. B174, 151, [JHEP 1986].

Mark Goodsell (2009), “Light Hidden U(1)‘s from String Theory,” in Proceedings, 5th Patras Workshop on Axions, WIMPs and WISPs (AXION-WIMP 2009): Durham, UK, July 13-17, 2009, pp. 165–168, arXiv:0912.4206 [hep-th].

Mark Goodsell, Joerg Jaeckel, Javier Redondo, and Andreas Ringwald (2009), “Naturally Light Hidden Photons in LARGE Volume String Compactifications,” JHEP 11, 027, arXiv:0909.0515 [hep-ph].

Ben-Ami Gradwohl, and Joshua A. Frieman (1992), “Dark matter, long range forces, and large scale structure,” Astrophys. J. 398, 407–424.

Peter W. Graham, Jeremy Mardon, and Surjeet Rajendran (2016), “Vector Dark Matter from Inflationary Fluctuations,” Phys. Rev. D 93 (10), 103520, arXiv:1504.02102 [hep-ph].

B. Grzadkowski, M. Ikrzyzynski, M. Misiak, and J. Rosiek (2010), “Dimension-Six Terms in the Standard Model Lagrangian,” JHEP 10, 085, arXiv:1008.4884 [hep-ph].

E. W. Hagley, and F. M. Pipkin (1994), “Separated oscillatory field measurement of hydrogen S-21/2- P-23/2 fine structure interval,” Phys. Rev. Lett. 72, 1172–1175.

Thomas Hambye, Michel H.G. Tytgat, Jérôme Vandecasteele, and Laurent Vanderheyden (2019), “Dark matter from dark photons: a taxonomy of dark matter production,” Phys. Rev. D 100 (9), 095018, arXiv:1908.0896 [hep-ph].

D. Hanneke, S. Fogwell, and G. Gabrielse (2008), “New Measurement of the Electron Magnetic Moment and the Fine Structure Constant,” Phys. Rev. Lett. 100, 120801, arXiv:0801.1134 [physics.atom-ph].

Edward Hardy, and Robert Lasenby (2017), “Stellar cooling bounds on new light particles: plasma mixing effects,” JHEP 02, 033, arXiv:1611.05852 [hep-ph].

Xiao-Gang He, Girish C. Joshi, H. Lew, and R. R. Volkas (1991), “Simplest Z-prime model,” Phys. Rev. D44, 2118–2132.
Jonathan J. Heckman, and Cumrun Vafa (2011), “An Exceptional Sector for F-theory GUTs,” Phys. Rev. D83, 026006, arXiv:1006.5459 [hep-th].

Julian Heeck (2014), “Unbroken B T L symmetry,” Phys. Lett. B759, 259–262, arXiv:1408.6845 [hep-ph].

A. Heister, et al. (ALEPH) (2003), “Single photon and multiphoton production in e+e− collisions at 3 GeV,” Eur. Phys. J. C28, 1–13.

J. L. Hewett, et al. (2012), “Fundamental Physics at the Intensity Frontier,” arXiv:1205.2671 [hep-ex].

Scott Hoffmann (1987), “Paraphotons and Axions: Similarities in Stellar Emission and Detection,” Phys. Lett. B193, 117–122.

Bob Holdom (1986a), “Searching for e Charges and a New U(1),” Phys. Lett. B178, 65–70.

Bob Holdom (1986b), “Two U(1)’s and Epsilon Charge Shifts,” Phys. Lett. B166B, 196–198.

Dan Hooper, Neal Weiner, and Wei Xue (2012), “Dark Forces and Light Dark Matter,” Phys. Rev. D86, 056009, arXiv:1206.2929 [hep-ph].

C. L. Hsu, et al. (Belle) (2012), “Search for B0 decays to invisible final states,” Phys. Rev. D86, 032002, arXiv:1206.5948 [hep-ex].

Philip Ilten, Yotam Soreq, Jesse Thaler, Mike Williams, and Wei Xue (2015), “Dark photons from charm mesons at LHCb,” Phys. Rev. D92 (11), 051105, arXiv:1409.6779 [hep-ex].

Philip Ilten, Yotam Soreq, Mike Williams, and Wei Xue (2018), “Serendipity in dark photon searches,” JHEP 06, 004, arXiv:1801.04847 [hep-ph].

Philip Ilten, Jesse Thaler, Mike Williams, and Wei Xue (2015), “Dark photons from charm mesons at LHCb,” Phys. Rev. D92 (11), 115017, arXiv:1505.00011 [hep-ph].

Joerg Jaeckel, Martin Jankowiak, and Michael Spannowsky (2013), “LHC probes the hidden sector,” Phys. Dark Univ. 2, 111–117, arXiv:1212.3620 [hep-ph].

Joerg Jaeckel, and Sabhyasachi Roy (2010), “Spectroscopy as a test of Coulomb’s law: A Probe of the hidden sector,” Phys. Rev. D82, 125020, arXiv:1008.3536 [hep-ph].

David E. Kaplan, Gordan Z. Krnjaic, Keith R. Rehermann, and Christopher M. Wells (2010), “Atomic Dark Matter,” JCAP 1005, 021, arXiv:0909.0753 [hep-ph].

Marek Karlinski, Matthew Low, Jonathan L. Rosner, and Lian-Tao Wang (2015), “Radiative return capabilities of a high-energy, high-luminosity e+e− collider,” Phys. Rev. D92 (3), 035010, arXiv:1503.07209 [hep-ph].

Wolfgang Keil, Hans-Thomas Janka, David N. Schramm, Gunter Sigl, Michael S. Turner, and John R. Ellis (1997), “A Fresh look at axions and SN-1987A,” Phys. Rev. D56, 2419–2432, arXiv:astro-ph/9612222 [astro-ph].

Kevin J. Kelly, and Yu-Dai Tsai (2019), “Proton fixed-target scintillation experiment to search for millicharged dark matter,” Phys. Rev. D100 (1), 015043, arXiv:1812.03998 [hep-ex].

V. Khachatryan, et al. (CMS) (2016), “A search for pair production of new light bosons decaying into muons,” Phys. Lett. B752, 146–168, arXiv:1506.00424 [hep-ex].

Simon Knapen, Tongyan Lin, and Kathryn M. Zurek (2017), “Light Dark Matter: Models and Constraints,” Phys. Rev. D96 (11), 115021, arXiv:1709.07882 [hep-ph].

Joachim Kopp, Ranjan Laha, Toby Oplerkuh, and William Shepherd (2018), “Cuckoo’s eggs in neutron stars: can LIGO hear chirps from the dark sector?” JHEP 11, 096, arXiv:1807.02527 [hep-ph].

Seth Koren, and Robert Mc Gee (2020), “Freezing-in twin dark matter,” Phys. Rev. D101 (5), 055024, arXiv:1908.03559 [hep-ph].

Ely D. Kovetz, Vivian Poulin, Vera Gluscevic, Kimberly K. Boddy, Rennan Barkana, and Marc Kamionkowski (2018), “Tighter limits on dark matter explanations of the anomalous EDGES 21 cm signal,” Phys. Rev. D98 (10), 103529, arXiv:1807.11482 [astro-ph.CO].

J. P. Lees, et al. (BaBar) (2012a), “Improved Limits on B0 Decays to Invisible Final States and to ννγ,” Phys. Rev. D86, 051105, arXiv:1206.2543 [hep-ex].

J. P. Lees, et al. (BaBar) (2014), “Search for a Dark Photon in e+e− Collisions at BaBar,” Phys. Rev. Lett. 113 (20), 201801, arXiv:1406.2980 [hep-ex].

J. P. Lees, et al. (BaBar) (2017), “Search for Invisible Decays of a Dark Photon Produced in e+e− Collisions at BaBar,” Phys. Rev. Lett. 119 (13), 131804, arXiv:1702.03327 [hep-ex].

J.P. Lees, et al. (BaBar) (2012b), “Precision Measurement of the B → X s γ Photon Energy Spectrum, Branching Fraction, and Direct CP Asymmetry ACP(B → X s + γ),” Phys. Rev. Lett. 109, 191801, arXiv:1207.2690 [hep-ex].

Sebastian Lien, Gianfranco Bertone, Francesca Calore, Roberto Ruiz de Austri, Tim M. P. Tait, Roberto Trottta, and Christoph Weniger (2016), “Effective field theory of dark matter: a global analysis,” JHEP 09, 077, arXiv:1603.05994 [hep-ph].

Tongyan Lin, Hai-Bo Yu, and Kathryn M. Zurek (2012), “On Symmetric and Asymmetric Light Dark Matter,” Phys. Rev. D85, 065033, arXiv:1111.0293 [hep-ph].

Hongwan Liu, Nadav Joseph Outmezguine, Diego Redigolo, and Tomer Volansky (2019), “Reviving Millicharged Dark Matter for 21-cm Cosmology,” Phys. Rev. D 100 (12), 123011, arXiv:1908.06986 [hep-ph].

Gabriel Magill, Ryan Pilestid, Maxim Pospelov, and Yu-Dai Tsai (2019), “Millicharged particles in neutrino experiments,” Phys. Rev. Lett. 122 (7), 071801, arXiv:1806.03310 [hep-ph].

L. Marsicano, M. Battaglieri, M. Bondi, C.D. R. Carvaljal, A. Celentano, M. De Napoli, R. De Vita, E. Nardi, M. Raggi, and P. Valente (2018), “Dark photon production through positron annihilation in beam-dump experiments,” Phys. Rev. D 98 (1), 015031, arXiv:1802.03794 [hep-ex].

Samuel D. McDermott, and Samuel J. Witte (2020), “Cosmological evolution of light dark matter,” Phys. Rev. D 101 (6), 063030, arXiv:1911.05086 [hep-ph].

H. Merkel, et al. (2014), “Search at the Mainz Micromontron for Light Massive Gauge Bosons Relevant for the Muon g-2 Anomaly,” Phys. Rev. Lett. 112 (22), 221802, arXiv:1404.5502 [hep-ex].

Marcelo M. Miller Bertolami, Brenda E. Melendez, Leandro G. Altmann, and Jordi Isern (2014), “Revisiting the axion bounds from the Galactic white dwarf luminosity function,” JCAP 1410 (10), 069, arXiv:1406.7712 [hep-ph].

Peter Minkowski (1977), “μ → e γ at a Rate of One Out of 109 Muon Decays?” Phys. Lett. B67, 421–428.

Alessandro Mirizzi, Javier Redondo, and Gunter Sigl (2009), “Microwave Background Constraints on Mixing of Photons
with Hidden Photons,” JCAP 03, 026, arXiv:0901.0014 [hep-ph].

M. Misiak, et al. (2007), “Estimate of B(R → X, γ) at O(α2),” Phys. Rev. Lett. 98, 022002, arXiv:hep-ph/0609232.

Rabindra N. Mohapatra, and Jogesh C. Pati (1975a), “Left-Right Gauge Symmetry and an Isoconjugate Model of CP Violation,” Phys. Rev. D 11, 566–571.

Rabindra N. Mohapatra, and Goran Senjanovic (1981), “Lepton Number Parameters of Early Galaxies,” Astrophys. J. 243, 127–132.

M. Misiak, et al. (2007), “Estimate of B(R → X, γ) at O(α2),” Phys. Rev. Lett. 98, 022002, arXiv:hep-ph/0609232.

Raul A. Monsalve, Bradley Greig, Judd D. Bowman, An-

Cristina Mondino, Maxim Pospelov, Joshua T. Ruderman, Wei-Tou Ni, Sheau-Shi Pan, Hsien-Chi Yeh, Li-Shing Hou, Yuichiro Nakai, Ryo Namba, and Ziwei Wang (2020), “Results and perspectives in dark photon physics,” Riv. Nuovo Cim. 38 (10), 449–505.

Mauro Raggi, and Venelin Kozhuharov (2015), “Results and perspectives in dark photon physics,” Riv. Nuovo Cim. 38 (10), 449–505.

Brian Patt, and Frank Wilczek (2006), “Higgs-field portal Dark Matter and the Misalignment Mechanism,” Phys. Rev. D 73, 105001, arXiv:1105.2812 [hep-ph].

Wei-Tou Ni, Sheau-Shi Pan, Hsien-Chi Yeh, Li-Shing Hou, and Ju-Ling Wan (1999), “Search for an axionlike spin coupling using a paramagnetic salt with a dc SQUID,” Phys. Rev. Lett. 82, 2439–2442.

L. B. Okun (1982), “LIMITS OF ELECTRODYNAMICS: PARAPHOTONS?” Sov. Phys. JETP 56, 502, [Zh. Eksp. Teor. Fiz.83,892(1982)].

Jogesh C. Pati, and Abdus Salam (1974), “Lepton Number as the Fourth Color,” Phys. Rev. D 10, 275–289, [Erratum: Phys.Rev.D 11, 703–703 (1975)].

Ann E. Nelson, and Jakub Scholtz (2011), “Dark Light, Dark Matter and the Misalignment Mechanism,” Phys. Rev. D 84, 103501, arXiv:1105.2812 [hep-ph].

D. Thomas (2015), “Axion Bremsstrahlung in Dense Stars,” Astrophys. J. 800, 449–505.

Annika H. G. Peter, Miguel Rocha, James S. Bullock, and Manoj Kaplinghat (2013), “Cosmological Simulations of Self-Interacting Dark Matter II: Halo Shapes vs. Observations,” Mon. Not. Roy. Astron. Soc. 430, 105, arXiv:1208.3026 [astro-ph.CO].

Maxim Pospelov (2009), “Secluded U(1) below the weak scale,” Phys. Rev. D 80, 095002, arXiv:0811.1030 [hep-ph].

Maxim Pospelov, Adam Ritz, and Mikhail B. Voloshin (2008), “Bosonic super-WIMPs as keV-scale dark matter,” Phys. Rev. D 78, 115012, arXiv:0807.3279 [hep-ph].

Maxim Pospelov, and Tommasi ter Veldhuis (2000), “Direct and indirect limits on the electromagnetic formfactors of WIMPs,” Phys. Lett. B 480, 181–186, arXiv:hep-ph/0003010.

John Preskill, Mark B. Wise, and Frank Wilczek (1983), “Cosmology of the Invisible Axion,” Phys. Lett. 120B, 127–132.

A. A. Prinz, et al. (1998), “Search for millicharged particles at SLAC,” Phys. Rev. Lett. 81, 1175–1178, arXiv:hep-ex/9804008 [hep-ex].

G. G. Raffelt (1996), Stars as laboratories for fundamental physics (Chicago, USA: Univ. Pr. (1996) 664 p).

Georg Raffelt, and David Seckel (1995), “A selfconsistent approach to neutral current processes in supernova cores,” Phys. Rev. D 52, 1780–1799, arXiv:astro-ph/9312019.

Georg G. Raffelt (1990), “Axion bremsstrahlung in red giants,” Phys. Rev. D 41, 1324–1326.
13 TeV,” JHEP 02, 074, arXiv:1810.00196 [hep-ex].
Albert M. Sirunyan, et al. (CMS) (2019c), “Search for supersymmetric partners of electrons and muons in proton-proton collisions at \( \sqrt{s} = 13 \) TeV,” Phys. Lett. B790, 140-166, arXiv:1806.05264 [hep-ex].
A. Sommerfeld (1931), “Über die beugung und bremsung der elektronen.” Annalen der Physik 403, 257.
Jhih-Ying Su, and Jusak Tandean (2019), “Searching for dark photons in hyperon decays,” arXiv:1911.13301 [hep-ph].
Jhih-Ying Su, and Jusak Tandean (2020), “Seeking massless dark photons in the decays of charmed hadrons,” arXiv:2005.05297 [hep-ph].
M. Tanabashi, et al. (Particle Data Group) (2018), “Review of Particle Physics,” Phys. Rev. D 98 (3), 030001.
Javier Tiffenberg, Miguel Sofo-Haro, Alex Drlica-Wagner, Rouven Essig, Yann Guardincerri, Steve Holland, Tomer Volansky, and Tien-Tien Yu (SENSEI) (2017), “Single-electron and single-photon sensitivity with a silicon Skipper CCD,” Phys. Rev. Lett. 119 (13), 131802, arXiv:1706.00028 [physics.ins-det].
Sean Tulin, Hai-Bo Yu, and Kathryn M. Zurek (2013), “Beyond Collisionless Dark Matter: Particle Physics Dynamics for Dark Matter Halo Structure,” Phys. Rev. D87 (11), 115007, arXiv:1302.3898 [hep-ph].
Nicolás Viaux, Márcio Catelan, Peter B. Stetson, Georg Raffelt, Javier Redondo, Aldo A. R. Valcarce, and Achim Weiss (2013), “Neutrino and axion bounds from the globular cluster M5 (NGC 5904),” Phys. Rev. Lett. 111, 231301, arXiv:1311.1669 [astro-ph.SR].
Hendrik Vogel, and Javier Redondo (2014), “Dark Radiation constraints on minicharged particles in models with a hidden photon,” JCAP 1402, 029, arXiv:1311.2600 [hep-ph].
Digvijay Wadekar, and Glennys R. Farrar (2019), “First direct astrophysical constraints on dark matter interactions with ordinary matter at very low velocities,” arXiv:1903.12190 [hep-ph].
D. J. Wineland, J. J. Bollinger, D. J. Heinzen, W. M. Itano, and M. G. Raizen (1991), “Search for anomalous spin-dependent forces using stored-ion spectroscopy,” Phys. Rev. Lett. 67, 1735–1738.
Samuel J. Witte, Salvador Rosauro-Alcaraz, Samuel D. McDermott, and Vivian Poulin (2020), “Dark Photon Dark Matter in the Presence of Inhomogeneous Structure,” arXiv:2003.13698 [astro-ph.CO].
Tsutomu Yanagida (1980), “Horizontal Symmetry and Masses of Neutrinos,” Prog. Theor. Phys. 64, 1103.