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Letter of interest:

The massless dark photon as a benchmark model

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Abstract

In this LOI, after briefly recalling the difference between the massive and the massless case, we explain why we think that the massless dark photon provides an interesting benchmark model in the search for a dark sector. We review the main constraints of the parameters of the model and discuss what are the most promising experiments and their discovery potential.

LET US BRIEFLY RECALL how the two kinds, massless and massive, of dark photons arise. The most general kinetic part of the Lagragian of two $U(1)_a$ and $U(1)_b$ gauge bosons is

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

The gauge boson A_b^μ is taken to couple to the current J_μ of ordinary SM matter, the other, A_a^μ, to the current J'_μ, which is made of dark-sector matter:

$$\mathcal{L} = e\, J_\mu A_b^\mu + e' J'_\mu A_a^\mu , \tag{2}$$

with e and e' the respective coupling constants.

As first discussed in [1], the classical Lagrangian can be diagonalized. What happens at the quantum level and how the mixing manifests itself has been analyzed in detail in [2] for the unbroken gauge theory as well as the spontaneously broken case (see, also, the appendix of [3]). We follow here the recent review [4].

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

$$\begin{pmatrix} A_a^\mu \\ A_b^\mu \end{pmatrix} = \begin{pmatrix} \dfrac{1}{\sqrt{1-\varepsilon^2}} & 0 \\ -\dfrac{\varepsilon}{\sqrt{1-\varepsilon^2}} & 1 \end{pmatrix} \begin{pmatrix} \cos\theta & -\sin\theta \\ \sin\theta & \cos\theta \end{pmatrix} \begin{pmatrix} A'^\mu \\ A^\mu \end{pmatrix} , \tag{3}$$

where now we can identify A^μ with the ordinary photon and A'^μ with the dark photon. The additional orthogonal rotation in Eq. (3) is always possible and introduces an angle θ which is arbitrary as long as the gauge bosons are massless.

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

$$\begin{aligned} \mathcal{L}' \;=\; & \left[\frac{e' \cos\theta}{\sqrt{1-\varepsilon^2}} J'_\mu + e \left(\sin\theta - \frac{\varepsilon \cos\theta}{\sqrt{1-\varepsilon^2}} \right) J_\mu \right] A'^\mu \\ + \; & \left[-\frac{e' \sin\theta}{\sqrt{1-\varepsilon^2}} J'_\mu + e \left(\cos\theta + \frac{\varepsilon \sin\theta}{\sqrt{1-\varepsilon^2}} \right) J_\mu \right] A^\mu . \end{aligned} \tag{4}$$

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

$$\boxed{\mathcal{L}' = \left[\frac{e'}{\sqrt{1-\varepsilon^2}} J'_\mu - \frac{e\varepsilon}{\sqrt{1-\varepsilon^2}} J_\mu \right] A'^\mu + e J_\mu A^\mu .} \tag{5}$$

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 $\varepsilon e/\sqrt{1-\varepsilon^2}$ 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 [4]. 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:

$$\boxed{\mathcal{L}' = e' J'_\mu A'^\mu + \left[-\frac{e'\varepsilon}{\sqrt{1-\varepsilon^2}} J'_\mu + \frac{e}{\sqrt{1-\varepsilon^2}} J_\mu \right] A^\mu} \tag{6}$$

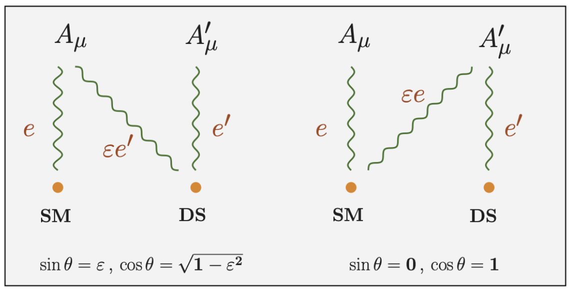

Figure 1: Scheme of the coupling of the ordinary (A_μ) and dark (A'_μ) photon to the SM and dark-sector (DS) particles for the two choices of the angle θ discussed in the main text. e and e' are the couplings of the ordinary and dark photons to their respective sectors.

If the gauge symmetry is spontaneously broken, the diagonalization of the mass terms locks the angle θ 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 [5] for a review and the relevant references)

$$\mathcal{L}_{Stu} = -\frac{1}{2} M_a^2 A_{a\mu} A_a^\mu - \frac{1}{2} M_b^2 A_{b\mu} A_b^\mu - M_a M_b A_{a\mu} A_b^\mu \,. \tag{7}$$

In this case, as in the spontaneously broken case, the angle θ is fixed and equal to

$$\sin\theta = \frac{\delta\sqrt{1-\varepsilon^2}}{\sqrt{1-2\delta\varepsilon+\delta^2}} \qquad \cos\theta = \frac{1-\delta\varepsilon}{\sqrt{1-2\delta\varepsilon+\delta^2}} \tag{8}$$

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

$$\begin{aligned} \mathcal{L}'' &= \frac{1}{\sqrt{1-2\delta\varepsilon+\delta^2}} \left[\frac{e'\,(1-\delta\varepsilon)}{\sqrt{1-\varepsilon^2}} J'_\mu + \frac{e\,(\delta-\varepsilon)}{\sqrt{1-\varepsilon^2}} J_\mu \right] A'^\mu \\ &+ \frac{1}{\sqrt{1-2\delta\varepsilon+\delta^2}} \left[e J_\mu - \delta e' J'_\mu \right] A^\mu \,. \end{aligned} \tag{9}$$

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 θ 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 the SM particles (in addition to that to dark-sector matter) and described by the Lagrangian

$$\boxed{\mathcal{L} \supset -\frac{e\varepsilon}{\sqrt{1-\varepsilon^2}} J_\mu A'^\mu \simeq -e\,\varepsilon\, J_\mu A'^\mu \,,} \tag{10}$$

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 $\tan\theta = \left[\varepsilon/\sqrt{1-\varepsilon^2}\right]$—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.

THE MASSLESS DARK PHOTON does not interact directly with the currents of the SM fermions, as shown by the Lagrangian in Eq. (6). The higher-order operators through which the interaction with ordinary matter ψ^i takes place start with the dimension-five operators in the Lagrangian

$$\mathcal{L} = \frac{e_D}{2\Lambda_5}\overline{\psi}^i\,\sigma_{\mu\nu}\left(\mathbb{D}_M^{ij} + i\gamma_5\,\mathbb{D}_E^{ij}\right)\psi^j\,F'^{\mu\nu}\,, \tag{11}$$

where $F'_{\mu\nu}$ is the field strength associated to the dark photon field A'_μ, and $\sigma_{\mu\nu} = i/2\,[\gamma_\mu, \gamma_\nu]$. The operator proportional to the coefficient $\mathbb{D}_M$ is the magnetic dipole moment and that proportional to the coefficient $\mathbb{D}_E$ 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. (11) 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 ψ_L and singlets ψ_R. In this case, the operators contain the Higgs boson field and can be written as

$$\boxed{\mathcal{L} = \frac{e_D}{2\Lambda^2}\overline{\Psi}_L^i\,\sigma_{\mu\nu}\left(\mathbb{D}_M^{ij} + i\gamma_5\,\mathbb{D}_E^{ij}\right)H\psi_R^j\,F'^{\mu\nu} + \text{H.c.}} \tag{12}$$

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.

The scale Λ depends on the parameters of the underlaying 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. (12), as they require a chirality flip, can turn out to be enhanced, or suppressed, according to the underlaying 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.

THE PHENOMENOLOGY OF THE MASSLESS DARK PHOTON depends on the effect of the higher-order operator in Eq. (12) which mediates its interaction with the SM particles. This operator enters the measured observables with an effective scale Λ and the absolute value

$$d_M^{ij} \equiv |\mathbb{D}_M^{ij}| \tag{13}$$

of the magnetic dipole coefficient (neglecting the CP-odd $\mathbb{D}_E$) 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 Λ, the dipole coefficient d_M^{ij} 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

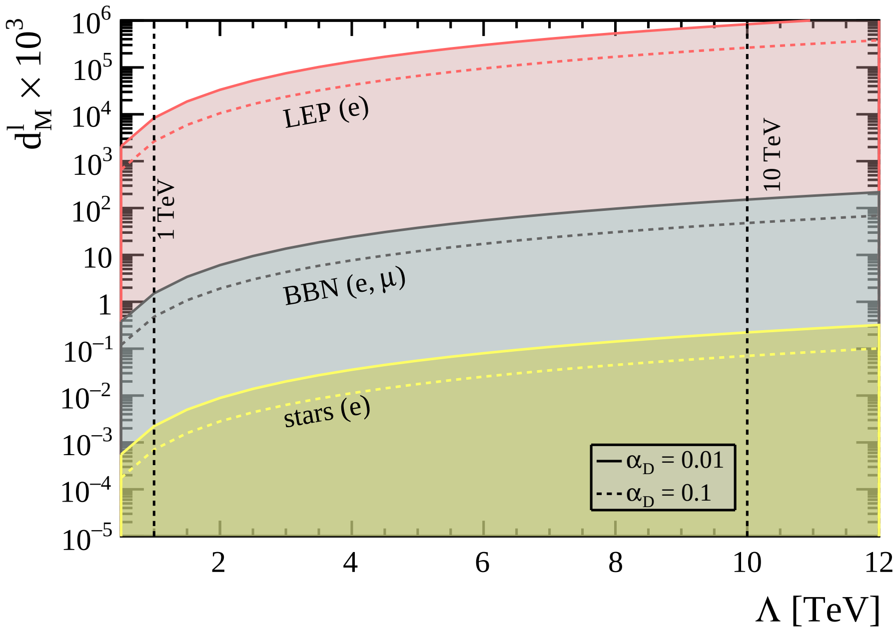

Figure 2: Model-independent limits for the interaction with **leptons**. The limits on the dark dipole operator d_M^ℓ/Λ^2 are shown by taking the coefficient d_M^ℓ as a function of the scale Λ (for two representative values of α_D). Given an energy scale, the allowed values for d_M^ℓ can be read from the plot. The strongest bound on electrons comes from stellar cooling (stars). Big bang nucleosynthesis (BBN) and collider physics (LEP) set the other depicted bounds. Solid lines are for the representative value $\alpha_D = 0.01$, dashed lines for $\alpha_D = 0.1$. [From ref [4]]

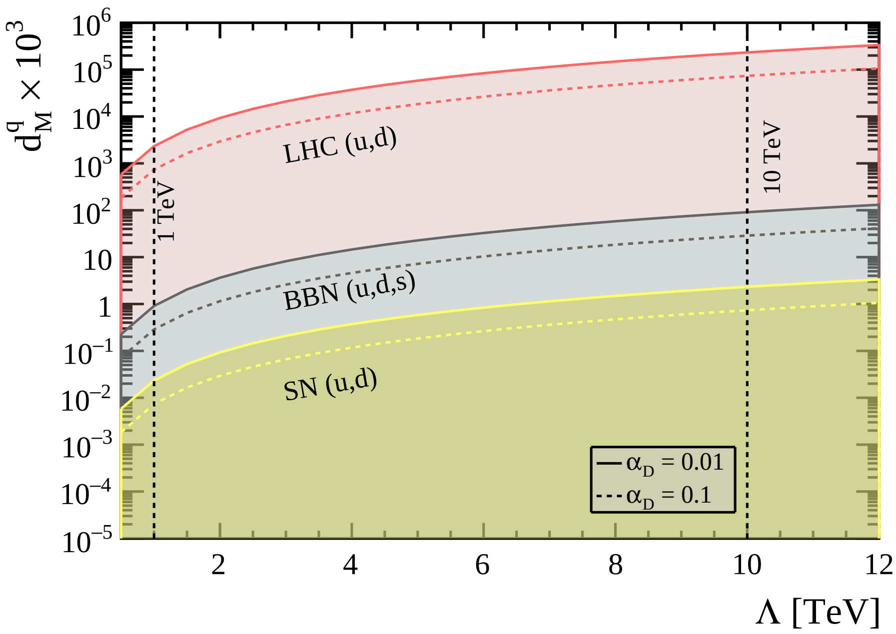

Figure 3: Model-independent limits for for the interaction with **quarks**. Same as in Fig. 2. The strongest bounds on light quarks comes from supernovae (SN). Primordial nucleosynthesis (BBN) and collider physics (LHC) set the other depicted bounds. Solid lines are for the representative value $\alpha_D = 0.01$, dashed lines for $\alpha_D = 0.1$. [From ref [4]]

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.

The known limits on losses of energy in stars and supernovae severely constrain the size of the parameter d_M/Λ^2. Further limits come from primordial nucleosynthesis. We show in Fig. 2 and 3 the more stringent limits. Though these limits are on the combinations d_M/Λ^2, with a factor depending on α_D, we find it convenient to plot them as d_M as a function of Λ so as to easily see what values of the dipole coefficient are allowed given a value for the scale Λ (and two representative value of α_D).

THE SIGNATURE OF THE MASSLESS DARK PHOTON can be searched in several experiments, some of which are already operating. We list here the most promising.

- 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 above applies given the flavor off-diagonal nature of the dipole operator in these cases.
  - **Processes in Kaon physics at NA62, NA64 and KOTO**: 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 [7] is within reach of the current sensitivity. The rare decays $K^+ \to \pi^+ \nu\bar{\nu}$ [8] and $K_L \to \pi^0 \nu\bar{\nu}$ [9] are other two processes where the physics of the dark photon can play a crucial role [10];
  - **Decays at BESIII**: Hyperion decays can be used for detecting the production of A' [11] and in the decay of charmed hadrons [12]'
  - **Decays into invisible states**: B-mesons at BaBar [13] and Belle [14] and $K_{L,S}$ and other neutral mesons at NA64 [15, 16] can be used to study the dark sector (assuming the invisible states belong to it). These decays are greatly enhanced by the Fermi-Sommerfeld [17, 18] effect due to their interaction with the massless dark photon—the same way as ordinary decays, like the β-decay, are enhanced by the same effect—making this another exciting area for searching the dark sector [19].
- Higgs and Z physics: The striking signature of a mono-photon plus missing energy can be used to search Higgs [20, 21, 22] and Z-boson [23, 24] decay into a visible and a dark photon. Again, the stringent astrophysical constrains discussed above 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.
- 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, 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) [25, 26]. 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 [27, 28, 29, 30].

To conclude, we think that the inclusion of a specific benchmark for the massless dark photon is important in making the search of a dark sector complete. The massless dark photon should not be relinquished into a generic study of higher order operators. It has its own characteristic features and provide a compelling physical model with many significative experimental signatures.

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