Hunting Down Massless Dark Photons in Kaon Physics

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If dark photons are massless, they couple to standard-model particles only via higher dimensional operators, while direct (renormalizable) interactions induced by kinetic mixing, which motivates most of the current experimental searches, are absent. We consider the effect of possible flavor-changing magnetic-dipole couplings of massless dark photons in kaon physics. In particular, we study the branching ratio for the process $K^+ \to \pi^+\pi^0\bar{\gamma}$ with a simplified-model approach, assuming the chiral quark model to evaluate the hadronic matrix element. Possible effects in the K^0-$\bar{K}^0$ mixing are taken into account. We find that branching ratios up to $O(10^{-7})$ are allowed—depending on the dark-sector masses and couplings. Such large branching ratios for $K^+ \to \pi^+\pi^0\bar{\gamma}$ could be of interest for experiments dedicated to rare K^+ decays like NA62 at CERN, where $\bar{\gamma}$ can be detected as a massless invisible system.

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The clarification of the origin of dark matter (DM) might require the existence of a *dark* sector made up of particles uncharged under the standard model (SM) gauge group. The possibility of extra secluded $U(1)$ gauge groups—mediating interactions in the dark sector via *dark* photons—is the subject of many experimental searches (see Ref. [1] for recent reviews). These searches are mostly based on the assumption that the secluded $U(1)$ gauge group is broken, and the corresponding *massive* dark photon (γ') interacts directly with the SM charged fields through renormalizable (dimension-four) operators induced by the kinetic mixing between dark and electromagnetic photons. Experimental results are then parametrized in terms of the dark-photon mass $m_{\gamma'}$ and mixing parameter ϵ, with dark photon signatures that can either correspond to its decay into SM particles or assume an invisible decay into extra dark fields. Because the induced operators have dimension four, most studies necessarily explore regions where the couplings are very small (*millicharges*).

We address instead the case of an unbroken dark $U(1)$ gauge symmetry, with a *massless* dark photon ($\bar{\gamma}$). The role of massless dark photons in galaxy formation and dynamics has been discussed in Refs. [2–6]. A strictly massless dark photon is very appealing from the theoretical point of view. Indeed, for massless dark photons it is possible [7] to define two fields, the dark and the ordinary photon, in such a manner that the dark photon only sees the dark sector. In this basis, ordinary photons couple to both the SM and the dark sector—the latter with millicharged strength to prevent macroscopic effects. Massless dark photons therefore interact with SM fields only through higher dimensional operators—typically suppressed by the mass scales related to new massive fields charged under the unbroken dark $U(1)$ gauge symmetry [8]—while their coupling constants can take natural values thanks to the built-in suppression associated to the higher dimensional operators. This makes the $\bar{\gamma}$ direct production in SM particle scattering or decay small and unobservable, consequently evading most of the search strategies for dark photons currently ongoing in laboratories. A possible exception is provided by the Higgs boson decay into dark photons in the nondecoupling regime. This scenario has been considered in Ref. [9], where observable $\bar{\gamma}$ production rates mediated by the Higgs decay $H \to \gamma\bar{\gamma}$ have been found at the LHC in realistic frameworks [10,11]. Flavor-changing-neutral-current (FCNC) decays of heavy flavors into a massless dark photon, $f \to f'\bar{\gamma}$, can offer other search channels with potentially observable rates [8,12].

Here we focus on FCNC effects induced by massless dark photons $\bar{\gamma}$ in kaon physics, and discuss the change of picture with respect to the massive case.

The kaon system can be studied with great accuracy, allowing us to probe indirectly energy scales as large as tens of TeV, hence, crucially constraining possible SM extensions. The detection of massive dark photons in K decays is presently under scrutiny [1,13]. One can consider *radiative* K decays where the (off-shell) SM photon γ is replaced by a γ', and look for resonances at $m_{\gamma'}$ for either e^+e^- ($\mu^+\mu^-$) final states, or (in case γ' decays into dark particles) for invisible final systems with a peak structure at $m_{\gamma'}$ in the missing mass distribution. Particular emphasis has been given to the decays $K^+ \to \pi^+\gamma'$ and $K^+ \to \mu^+\nu\gamma'$ [14–18]. However, if the secluded $U(1)$ gauge group is unbroken, these two channels are not viable. Indeed, $K^+ \to \pi^+\bar{\gamma}$ violates angular momentum conservation, while $K^+ \to \mu^+\nu\bar{\gamma}$ would require unsuppressed $\bar{\gamma}$ couplings.

Because K^+ decays into a dark photon $\bar{\gamma}$ must necessarily proceed through short-distance effects, we argue that the most interesting channel to look for massless dark photons in kaon physics could be the decay $K^+ \to \pi^+\pi^0\bar{\gamma}$. This decay can be mediated by the FCNC transition $s \to d\bar{\gamma}$, prompted by a magnetic-dipole-type coupling generated at one loop by the dark-sector degrees of freedom. The dark photon gives rise in this case to a massless missing-momentum system inside the final state. Recently, the sensitivity of the NA62 experiment at the CERN SPS [19] to two-body K decays into a light vector decaying invisibly $[K^+ \to \pi^+ + (\gamma' \to E_{\rm miss})]$ has been emphasized [13]. For the three-body $K^+ \to \pi^+\pi^0\bar{\gamma}$ channel, whose kinematics is less characterized, the detection efficiency is expected to be less favorable. Nevertheless—since the $K^+ \to \pi^+\pi^0\bar{\gamma}$ channel has a unique potential to unveil the existence of a massless dark photon—we think that the NA62 Collaboration should consider search strategies aiming at detecting this newly proposed process, whose branching ratio (BR) can reach 10^{-7} in a *simplified* model of the dark sector, as we estimate in the following.

A simplified model of the dark sector.—We estimate $\mathrm{BR}(K^+ \to \pi^+\pi^0\bar{\gamma})$ in a simplified model that makes as few assumptions as possible, while providing the dipole-type transition we are interested in.

The minimal choice in terms of fields consists of a SM extension where there is a new (heavy) dark fermion Q, singlet under the SM gauge interactions, but charged under an unbroken $U(1)_D$ gauge group associated to the massless dark photon. SM fermions couple to the dark fermion by means of a Yukawa-like interaction in the Lagrangian $\mathcal{L}$,

$$\mathcal{L} \supset g_L(\bar{Q}_L q_R)S_R + g_R(\bar{Q}_R q_L)S_L + \mathrm{H.c.}, \tag{1}$$

where new (heavy) *messenger* scalar particles, S_L and S_R, enter as well. In Eq. (1), the q_L and q_R fields are the SM fermions [$SU(3)$ triplets and, respectively, $SU(2)$ doublets and singlets]. Flavor indices are implicit, and we assume common (i.e., flavor blind) couplings g_L and g_R. The left-handed messenger field S_L is a $SU(2)$ doublet, the right-handed messenger field S_R is a $SU(2)$ singlet, and both are $SU(3)$ color triplets. These messenger fields are charged under $U(1)_D$, carrying the same $U(1)_D$ charge of the dark fermion.

In order to generate chirality-changing processes we also need in the Lagrangian the mixing terms

$$\mathcal{L} \supset \lambda_S S_0(S_L S_R^\dagger \tilde{H}^\dagger + S_L^\dagger S_R H), \tag{2}$$

where H is the SM Higgs boson, $\tilde{H} = i\sigma_2 H^\star$, and S_0 a scalar singlet. The Lagrangian in Eq. (2) gives rise to the mixing after both the S_0 and H scalars take a vacuum expectation value (VEV), respectively, μ_S and v—the electroweak VEV. After diagonalization, the messenger fields $S_\pm$ couple to both left- and right-handed SM fermions with strength $g_L/\sqrt{2}$ and $g_R/\sqrt{2}$, respectively. We can assume that the size of this mixing—proportional to the product of the VEVs ($\mu_s v$)—is large and of the same order of the masses of the heavy fermion and scalars.

The SM Lagrangian plus the terms in Eqs. (1)–(2) (supplemented by the corresponding kinetic terms) provide a simplified model for the dark sector and the effective interaction of the SM degrees of freedom with the massless dark photon $\bar{\gamma}$. SM fermions couple to $\bar{\gamma}$ only via nonrenormalizable interactions, induced by loops of the dark-sector states. Two scales are relevant: the dark fermion mass M_Q, which parametrizes the chiral symmetry breaking in the dark sector, and the lightest-messenger mass scale m_S. Since we are considering the contribution to the magnetic dipole operator (assuming vanishing quark masses), the dominant effective scale associated with it will either be chirally suppressed (being proportional to M_Q/m_S^2, for $m_S \gg M_Q$), or scale as $1/M_Q$ (for $m_S \ll M_Q$) due to decoupling. In order to have only one dimensionful parameter, in our analysis we assume a common mass for the dark fermion and the lightest scalar field, which we identify with the new-physics scale Λ. This choice corresponds to the maximum chiral enhancement.

This scenario is a simplified version of the model in Refs. [10–12] (possibly providing a natural solution to the SM flavor-hierarchy problem), as well as a template for many models of the dark sector.

Bounds from $K^0 - \bar{K}^0$ and astrophysics.—A most stringent limit to the mass scale and couplings of the above simplified model comes from its extra contributions to the K^0-$\bar{K}^0$ mixing in the kaon system (related to the mass difference ΔM_K of the neutral mass eigenstates K_L and K_S, assuming CPT).

In order to compute the dark-sector effects on ΔM_K, we need to evaluate the dark-sector contribution to the effective Hamiltonian for the $\Delta S = 2$ transitions, $\mathcal{H}_{\rm eff}^{\Delta S=2}$

$$\Delta M_K = 2\mathrm{Re}[\langle K^0|\mathcal{H}_{\rm eff}^{\Delta S=2}|\bar{K}^0\rangle]. \tag{3}$$

The scalar-fermion interaction in Eq. (1) induces a new set of operators, which are reported in Table I, then obtaining

$$\mathcal{H}_{\rm eff}^{\Delta S=2} = \sum_i^5 C_i Q_i + \sum_{i=1}^3 \tilde{C}_i \tilde{Q}_i. \tag{4}$$

The Wilson coefficients at the matching scale are computed by considering the exchange of the lightest messenger state in the loop, which provides a good estimate of the dominant contribution in the large-mixing limit of the messenger mass sector.

We compute the corresponding Wilson coefficients $C_i(\mu)$ at the $\mathcal{O}(\alpha_s)$ next-to-leading order, after running them from the matching scale down to the low energy scale $\mu \sim 2$ GeV, where the corresponding matrix elements are estimated on the lattice [21]. We assume as matching scale the characteristic mass Λ of the lightest-messenger and

TABLE I. In the first two rows, relevant operators are numbered according to the notation in [20,21]. The matrix elements $\langle K^0|Q_i|\bar{K}^0\rangle$ (in the vacuum insertion approximation for the renormalized operators Q_i at the low energy scale $\mu = 2$ GeV) are given in the third row multiplied by the respective bag factors $B_i(\mu)$ [21] evaluated at same scale, with $X_K(\mu) = (m_K/(m_d(\mu) + m_s(\mu)))^2$. The fourth row gives the Wilson coefficients at the matching scale (the common factor at the matching being $C^2 = \xi^2/(16\pi^2\Lambda^2)$ [9], where $\xi = g_L g_R/2$). Following [21], we take $m_d(\mu) = 7$ MeV, $m_s(\mu) = 125$ MeV, $m_K = 497$ MeV, $f_K = 160$ MeV, and $B_{1,2,3,4,5}(\mu) = 0.60, 0.66, 1.05, 1.03, 0.73$, respectively.

Q_1, $\tilde{Q}_1$	Q_2, $\tilde{Q}_2$	Q_3, $\tilde{Q}_3$	Q_4	Q_5
$\bar{d}^\alpha_L\gamma_\mu s^\alpha_L\bar{d}^\beta_L\gamma_\mu s^\beta_L$, $(L \leftrightarrow R)$	$\bar{d}^\alpha_R s^\alpha_L\bar{d}^\beta_R s^\beta_L$, $(L \leftrightarrow R)$	$\bar{d}^\alpha_R s^\beta_L\bar{d}^\beta_R s^\alpha_L$, $(L \leftrightarrow R)$	$\bar{d}^\alpha_R s^\alpha_L\bar{d}^\beta_L s^\beta_R$	$\bar{d}^\alpha_R s^\beta_L\bar{d}^\beta_L s^\alpha_R$
$1/3 m_K f_K^2 B_1(\mu)$	$-5/2 X_K m_K f_K^2 B_3(\mu)$	$1/24 X_K m_K f_K^2 B_3(\mu)$	$1/4 X_K m_K f_K^2 B_4(\mu)$	$1/12 X_K m_K f_K^2 B_5(\mu)$
$-1/24C^2$	0	$1/12C^2$	$1/6C^2$	$1/6C^2$

dark-fermion states, assumed to be equal. Following this procedure, the dark-sector contribution to ΔM_K (in TeV) is

$$\Delta M_K = 8.47 \times 10^{-13}\frac{\xi^2}{\Lambda^2}, \tag{5}$$

where $\xi = g_L g_R/2$, and Λ is in TeV units. We then assume that the above contribution of the new operators to Eq. (3) does not exceed 30% of the measured ΔM_K value [22]. Equation (5) turns then into an upper bound for the allowed values for the ξ^2/Λ^2 ratio.

While the flavor-changing dipole operator induced in the simplified model [see Eq. (6) below] *per se* is only bounded by kaon physics, if we make the (very conservative) assumption that the model also gives flavor-diagonal dipole operators and these are the same size in the quark and lepton sectors, a bound can be derived from stellar cooling carried out by the emission of massless dark photons. Under these assumptions, the lower limit on the energy scale Λ/ξ from K^0-$\bar{K}^0$ mixing in Eq. (5) falls between the current astrophysical bounds [23]—with the most stringent one from white-dwarf stars being 1 order of magnitude stronger and that from the Sun 1 order of magnitude weaker.

Amplitude and decay rate.—The $K^+ \to \pi^+\pi^0\bar{\gamma}$ decay originates from the dimension-five magnetic dipole operator $\hat{Q} = (\bar{s}\sigma^{\mu\nu}d)\bar{F}_{\mu\nu}$, where $\bar{F}_{\mu\nu}$ is the $\bar{\gamma}$ field strength, $\sigma_{\mu\nu} = \frac{1}{2}[\gamma_\mu, \gamma_\nu]$, and color and spin contractions are understood. $\hat{Q}$ enters the effective Hamiltonian for $\Delta S = 1$ transitions as

$$\mathcal{H}^{\Delta S=1}_{\rm eff} = \frac{e_D}{64\pi^2}\frac{\xi}{\Lambda}\hat{Q}, \tag{6}$$

where $\alpha_D = e_D^2/(4\pi)$ is the $\bar{\gamma}$ coupling strength. The Wilson coefficient multiplying the magnetic operator in Eq. (6) is obtained by integrating the vertex function in our simplified model (see Fig. 1). We have checked Eq. (6) by means of `Package X` [24].

The operator in Eq. (6) contributes only to the magnetic component of the process

$$K^+(p) \to \pi^+(q_1)\pi^0(q_2)\bar{\gamma}(k), \tag{7}$$

while its contribution to the process $K^+ \to \pi^+\bar{\gamma}$ identically vanishes. The amplitude $\hat{M} \equiv \langle\bar{\gamma}\pi^+\pi^0|\mathcal{H}^{\Delta S=1}_{\rm eff}|K^+\rangle$ in the momentum space can be written as

$$\hat{M} = \frac{M(z_1, z_2)}{m_K^3}\varepsilon_{\mu\nu\rho\sigma}q_1^\nu q_2^\rho k^\sigma \varepsilon^\mu(k), \tag{8}$$

where $\varepsilon^\mu(k)$ is the $\bar{\gamma}$ polarization vector. The corresponding differential decay rate is

$$\frac{d^2\Gamma}{dz_1 dz_2} = \frac{m_K}{(4\pi)^3}|M(z_1, z_2)|^2\{z_1 z_2[1 - 2(z_1 + z_2) - r_1^2 - r_2^2] - r_1^2 z_2^2 - r_2^2 z_1^2\}, \tag{9}$$

where $z_i = k \cdot q_i/m_K^2$ and $r_i = M_{\pi_i}/m_K$ [25].

The matrix element in Eq. (8) can be estimated by means of the chiral quark model (χQM) [26]. In this model quarks are coupled to hadrons by an effective interaction so that matrix elements can be evaluated by loop diagrams (see Fig. 2). In general, there are several free parameters, but in the present case only M, the mass of the constituent quarks, and f, the pion decay constant, enter the computation. The model has been applied to kaon physics in Ref. [27], where a fit of the CP preserving amplitudes of the nonleptonic

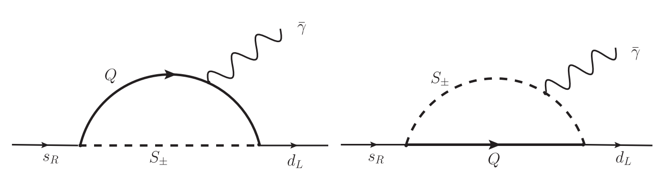

FIG. 1. Vertex diagrams for the generation of the dipole operator in the simplified model of the dark sector (same for the specific model in Refs. [10–12]).

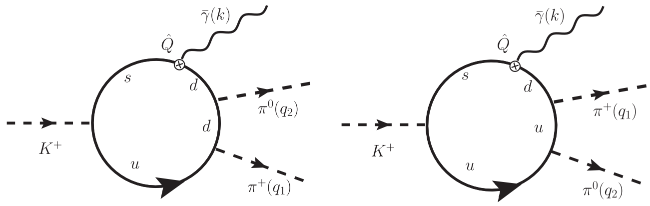

FIG. 2. χQM diagrams for the process $K^+ \to \pi^+\pi^0\tilde{\gamma}$. The crossed circle stands for the insertion of the magnetic dipole operator $\hat{Q}$ in Eq. (6).

decay of neutral kaons has yielded a value $M = 200$ MeV [28] with an error of less of 5%.

According to the χQM we obtain that the magnetic component generated by the dipole operator in Eq. (6) is given by

$$\begin{aligned}\frac{M(z_1,z_2)}{m_K^3} = {} & \frac{e_D}{32\pi^2}\frac{\xi}{\Lambda}\frac{M^3}{\pi^2 f^3}[M^2 D_0(0, m_\pi^2, m_\pi^2, m_K^2; 2m_K^2 z_1 \\ & + m_\pi^2, m_K^2(1-2z_1-2z_2); M, M, M, M). \\ & - .D_{00}(0, m_\pi^2, m_\pi^2, m_K^2; 2m_K^2 z_1 \\ & + m_\pi^2, m_K^2(1-2z_1-2z_2); \\ & \times M, M, M, M) + (z_1 \leftrightarrow z_2)], \end{aligned} \tag{10}$$

where D_0 and D_{00} are four-point Passarino-Veltman coefficient functions (see Ref. [29] for their explicit form) to be evaluated numerically [24].

Inserting the amplitude in Eq. (10) in the differential decay rate in Eq. (9) yields, after integration and by normalizing Γ by the total K^+ width $\Gamma_{\rm tot} = 5.317 \times 10^{-14}$ MeV [22],

$$\mathrm{BR}(K^+ \to \pi^+\pi^0\bar{\gamma}) \simeq 1.31\alpha_D \eta^2 \frac{\xi^2}{\Lambda^2}, \tag{11}$$

where we assumed $M = 200\,\mathrm{MeV}$, $f = 92.4\,\mathrm{MeV}$, $m_K = 494\,\mathrm{MeV}$, and $m_{\pi^+} = m_{\pi^0} = 136$ MeV. The coefficient η accounts for the renormalization of the Wilson coefficient of the dipole operator in going from the Λ scale to approximately m_K. We assume it equal to 1, and discuss the impact of possible uncertainties below.

$\mathrm{BR}(K^+ \to \pi^+\pi^0\bar{\gamma})$ is proportional to ξ^2/Λ^2, just as ΔM_K in Eq. (5). By taking for ξ^2/Λ^2 the value that saturates the ΔM_K constraint, we find an upper bound for the BR which is, for the representative value $\alpha_D = 0.1$,

$$\mathrm{BR}(K^+ \to \pi^+\pi^0\bar{\gamma}) \lesssim 1.6 \times 10^{-7}. \tag{12}$$

Figure 3 shows the $\mathrm{BR}(K^+ \to \pi^+\pi^0\bar{\gamma})$ contour plot versus the scale Λ and the coupling ξ, for $\alpha_D = 0.1$. We see that a rather large range of parameters is allowed for which the BR is sizable. The upper bound—given by Eq. (12)—is represented in Fig. 3 by the boundary of the gray area.

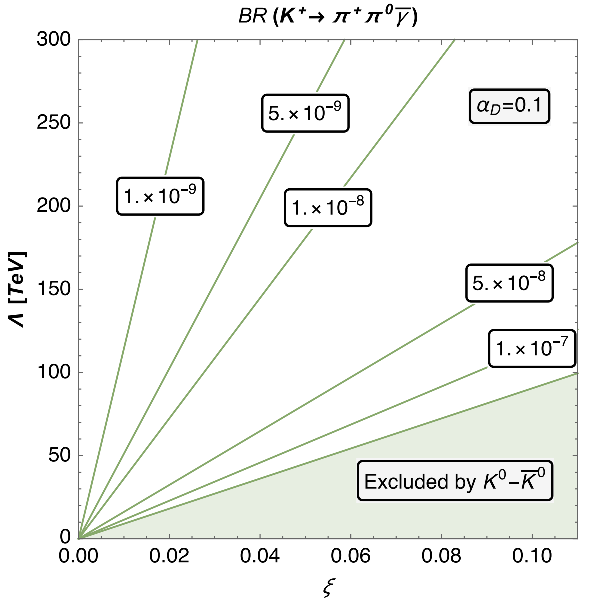

FIG. 3. $\mathrm{BR}(K^+ \to \pi^+\pi^0\bar{\gamma}))$ as a function of the effective scale Λ and coupling $\xi = g_L g_R/2$, for a representative choice of the coupling strength $\alpha_D = 0.1$.

There are three main sources of uncertainties in the result in Eq. (12): (i) The matrix element estimate computed in the χQM depends on the parameter M. The result in Ref. [28] seems to indicate a rather small uncertainty on this parameter but one must be aware of the dependence. We find an increase by a factor 2.5 in the BR when going from $M = 200$ to 250 MeV; (ii) Even though there are $O(p^4)$ chiral perturbation theory corrections to $K^+ \to \pi^+\pi^0\bar{\gamma}$, these have been shown to be small [30]; (iii) By taking the QCD leading-order multiplicative value $\eta = 0.5$ (at $\mu = 2$ GeV) [31], we find a BR smaller by a factor of 1/4. However, it is known that nonmultiplicative corrections go the opposite direction, and we thus need the (not yet available) complete evolution before trusting this correction. Moreover, the QCD renormalization introduces a strong dependence on the low-energy scale μ, because the matrix element computed within the χQM is scale independent. On top of these uncertainties, we have the overall dependence on the α_D strength on which the BR depends linearly. There exist cosmological relic density bounds on the ratio α_D/Λ^2 [3]. Our choice of $\alpha_D = 0.1$ is then consistent with Λ of the order of 10 TeV.

Similar predictions can be obtained in the specific flavor model of Refs. [10–12]. In particular, for $\alpha_D = 0.1$, the approximate upper bound is given by $\mathrm{BR} \simeq 1.2 \times 10^{-8}$. The lower BR is explained by the dark-fermion masses being related in this case to the radiative generation of SM Yukawa couplings, resulting in a stronger chiral suppression of the effective scale associated with the dipole operator $\hat{Q}$, which turns out to be proportional to the bottom-quark Yukawa coupling [12].

Conclusions.—NA62 at the CERN SPS will soon provide a sample of 10^{13} K^+, with hermetic photon coverage and good missing-mass resolution [19]. We propose to look for the rare decay $K^+ \to \pi^+\pi^0\bar{\gamma}$ (where $\bar{\gamma}$ gives rise to a massless invisible system) as a sensitive probe for massless dark photons, for which the presently most explored dark-photon channels mediated by kinetic-mixing interactions in kaon decays are nonviable.

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