Seeking massless dark photons in the decays of charmed hadrons

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Abstract

We entertain the possibility that massless dark photons exist and have nonnegligible flavor-changing dipole-type couplings to the $u$ and $c$ quarks. Such interactions give rise to the flavor-changing neutral current decays of charmed hadrons into a lighter hadron plus missing energy carried away by the dark photon. We investigate specifically the decays of the charmed mesons $D^+$, $D^0$, and $D_s^+$, bottom, charmed meson $B_c^+$, singly charmed baryons $\Lambda_c^+$, $\Xi_c^+$, and $\Xi_c^0$, and doubly-charmed baryon $\Xi_{cc}^{++}$. Employing a simplified new-physics model satisfying the relevant constraints, we find that these processes can have branching fractions reaching the $10^{-4}$ level. This suggests that one or more of these modes may be accessible in the near future by the ongoing Belle II and BESIII experiments. Since the same underlying operators are responsible for all of these transitions, detecting one of them automatically implies particular predictions for the others, allowing for additional experimental checks on the massless dark photon scenario.
I. INTRODUCTION

Attempts to address longstanding open questions in physics, such as the nature and origin of neutrino mass and the particle identity of cosmic dark matter, have increasingly postulated the presence of a dark sector beyond the standard model (SM). It is reasonable to expect that the new dark sector not only provides resolutions to some of those major puzzles but also furnishes extra ingredients which will facilitate further experimental access to it. Among the most attractive ones is a dark Abelian gauge group, $U(1)_D$, under which all SM fields are singlets. This symmetry may be spontaneously broken or stay unbroken, causing the associated gauge boson, the dark photon, to gain mass or remain massless.

Whether it is massive or massless, the hope is that the dark photon can somehow communicate with the SM as well as connect it to other constituents of the dark side, leading to interesting and potentially observable consequences. These possibilities have received a great deal of theoretical attention in the past few decades [1–20] and motivated numerous dedicated searches for dark photons [18–27], albeit still with negative outcomes to date. Most of these efforts have focused on the massive dark photon, $A'$, which can couple directly to SM fermions via the renormalizable operator $\epsilon e A'_\mu J_{\text{EM}}^\mu$ involving the electromagnetic current $e J_{\text{EM}}$ and a small constant $\epsilon$ brought about by the kinetic mixing between the dark and SM Abelian gauge fields [16–20]. In the presence of this coupling, $A'$ could be produced in the scattering or decays of leptons and quarks, including those of hadrons, and it could decay into electrically charged fermions or mesons. Because of these properties of $A'$, the measurements looking for it have been able to come up with limits on $\epsilon$ over various ranges of its mass [17–27].

If the dark photon is massless, the situation is very different but no less interesting. In this case, one can always define a linear combination of the dark and SM $U(1)$ gauge fields that has no renormalizable interactions with the SM and is identified with the massless dark photon [3, 6], which we denote by $\gamma$. The implications are that it has no direct couplings to SM fermions, in contrast to its massive counterpart, and that therefore restrictions inferred from the aforesaid quests for $A'$ do not apply to $\gamma$. However, the latter can still affect the SM sector through higher-dimensional operators induced by loop diagrams containing particles charged under $U(1)_D$ and also coupled to SM fields [6, 7]. This means that there may be more efficient ways to probe $\gamma$, which are worth pursuing and some of which will be explored in what follows.

In the absence of other particles beyond the SM lighter than the electroweak scale, the effective interactions of the massless dark photon with SM members can be described by operators which respect the SM gauge group and the unbroken $U(1)_D$. At leading order the couplings of $\gamma$ to quarks are of dipole type and given by the gauge-invariant Lagrangian [6]

$$\mathcal{L}_{\gamma q} = \frac{1}{\Lambda_{\gamma q}^2} \left( C_{jk} \overline{q}_j \sigma^{\mu\nu} d_k H + C'_{jk} \overline{u}_j \sigma^{\mu\nu} u_k \tilde{H} + \text{H.c.} \right) F_{\mu\nu},$$

where $\Lambda_{\gamma q}$ denotes an effective heavy mass, the $C$s are dimensionless coefficients which are generally
complex, \( q_j \) and \( d_k \) \((u_k)\) represent a left-handed quark doublet and right-handed down-type \((u\)-type) quark singlet, respectively, under the SU(2)\(_L\) gauge group, \( H \) stands for the SM Higgs doublet, \( \tilde{H} = i\tau_2 H^* \) with \( \tau_2 \) being the second Pauli matrix, \( \tilde{F}_{\mu\nu} = \partial_\mu \tilde{A}_\nu - \partial_\nu \tilde{A}_\mu \) is the dark photon’s field-strength tensor, \( \sigma^{\mu\nu} = i[\gamma^\mu, \gamma^\nu]/2 \), and summation over family indices \( j, k = 1, 2, 3 \) is implicit. Both \( \Lambda_{np} \) and the \( C_i \) depend on the details of the underlying new physics (NP), and in general \( C_{jk} \) and \( C'_{jk} \) are not necessarily related to one another.

Here we concern ourselves with the flavor-changing neutral current (FCNC) transitions among the lightest quarks which can arise from the operators in Eq. (1) and contribute to the decays of hadrons with missing energy. In the SM the corresponding processes involve an undetected neutrino pair in the final state and are highly suppressed, as they proceed from loop diagrams and are subject to the Glashow-Iliopoulos-Maiani (GIM) mechanism \[28, 29\]. It follows that these FCNC hadron decays with missing energy are potentially good places to seek signs of NP. If \( \mathcal{L}_{np} \) could impact these processes substantially, their rates might be amplified well above their SM expectations to values that are within the sensitivity reaches of ongoing high-intensity flavor experiments, such as Belle II \[30\], BESIII \[31, 32\], and NA62 \[33\].

The possibility that the massless dark photon, \( \tilde{\gamma} \), has flavor-changing interactions with the \( d \) and \( s \) quarks, via the \( \mathcal{C}_{jk} \) portion of \( \mathcal{L}_{np} \), has already been entertained to some extent previously. Such couplings generate the FCNC decays of kaons \[14\] and hyperons \[15\] with missing energy carried away by \( \tilde{\gamma} \). As pointed out in Refs. \[14, 15\], their respective branching fractions are permitted by current constraints to increase to levels which may be discoverable in the near future by NA62 and BESIII, respectively.

In this paper we concentrate on the case where \( \tilde{\gamma} \) has flavor-changing interactions with the lightest up-type quarks, \( u \) and \( c \), originating from the \( \mathcal{C}_{jk} \) parts in Eq. (1). After electroweak symmetry breaking, in the up-type-quark mass basis we can express the pertinent terms as

\[
\mathcal{L}_{uc\tilde{\gamma}} = \overline{u}(C + \gamma_5 C_5)\sigma^{\mu\nu}c\tilde{F}_{\mu\nu} + \text{H.c.},
\]

and so \( C = \Lambda_{np}^2 (C_{12} + C_{21}^*) v/\sqrt{8} \) and \( C_5 = \Lambda_{np}^2 (C_{12} - C_{21}^*) v/\sqrt{8} \) are parameters which have the dimension of inverse mass and are determined by the specifics of the ultraviolet-complete model, with \( v \approx 246 \text{ GeV} \) being the Higgs vacuum expectation value. In the next few sections, we analyze the implications of \( \mathcal{L}_{uc\tilde{\gamma}} \) for the FCNC decays of several charmed hadrons into two-body final states each consisting of a lighter hadron and \( \tilde{\gamma} \). These transitions have virtually no SM background because the corresponding SM contributions are the three-body modes with a neutrino pair which have tiny rates \[29\]. On the other hand, \( \mathcal{L}_{uc\tilde{\gamma}} \) could produce significantly enhancing effects on the two-body decays with the massless dark photon. Therefore, it is hoped that our study will motivate experimental efforts to pursue them, which may shed light on the existence of \( \tilde{\gamma} \).

We first deal with the decays of the lightest charmed-hadrons in Sec. II, namely the pseudoscalar mesons \( D^+ \), \( D^0 \), and \( D_s^+ \). For good measure, we include \( D^0 \to \gamma\tilde{\gamma} \), which is also caused by \( \mathcal{L}_{uc\tilde{\gamma}} \) and emits an ordinary photon, \( \gamma \), instead of a meson. Subsequently, we address the decay of the
bottom, charmed pseudoscalar-meson $B_c^+$ in Sec. [III]. Then we turn to the decays of singly charmed baryons $\Lambda_c^+, \Xi_c^+$, and $\Xi_c^0$ in Sec. [IV] and of the doubly-charmed baryon $\Xi_{cc}^{++}$ in Sec. [V]. Finally, we give our conclusions in Sec. [VI]. An appendix supplies extra information on the baryonic matrix elements of the dipole operators.

II. DECAYS OF CHARMED MESONS

The operators in $\mathcal{L}_{uc\gamma}$ can give rise to the FCNC decays of the lightest charmed-mesons into another meson and the dark photon. Given that the decay of a spinless particle into another spinless particle plus a massless gauge boson is forbidden by angular-momentum conservation and gauge invariance, in most of this section we consider modes where the initial particle is the charmed pseudoscalar-meson $D^+$ or $D^0$ or the charmed, strange pseudoscalar-meson $D_s^+$, whereas the daughter particles comprise a charmless vector-meson and $\gamma$. In particular, we look at $D^+ \rightarrow \rho^+\bar{\gamma}$, $D^0 \rightarrow \rho^0\bar{\gamma}$, $D^0 \rightarrow \omega\bar{\gamma}$, and $D_s^+ \rightarrow K^+\bar{\gamma}$. As for $D^0 \rightarrow \gamma\bar{\gamma}$, we will treat it towards the end of the section.

For $D^+ \rightarrow \rho^+\bar{\gamma}$ the amplitude $\mathcal{M}_{D^+\rightarrow\rho^+\bar{\gamma}}$ contains the mesonic matrix elements $\langle \rho^+ | \bar{u} \sigma^{\mu\nu} c | D^+ \rangle$ and $\langle \rho^+ | \bar{u} \sigma^{\mu\nu} \gamma_5 c | D^+ \rangle$, the general formulas for which are well known in the literature [34–36]. In $\mathcal{M}_{D^+\rightarrow\rho^+\bar{\gamma}}$ these matrix elements are to be contracted with the outgoing dark photon’s momentum $\vec{q}$ and polarization vector $\vec{\varepsilon}$. These two four-vectors fulfill the gauge requirement $\vec{\varepsilon} \cdot \vec{q} = 0$ and the masslessness condition $q^2 = 0$ due to $\bar{\gamma}$ being on-shell. Accordingly, we can write

$$\langle \rho^+(k) | \bar{u} \sigma^{\mu\nu} c | D^+(k+\bar{q}) \rangle \varepsilon^*_\mu \bar{q}_\nu = 2i f_{D^+\rho^+} \varepsilon^{\mu\nu} \varepsilon^*_\mu \bar{q}_\nu,$$

$$\langle \rho^+(k) | \bar{u} \sigma^{\mu\nu} \gamma_5 c | D^+(k+\bar{q}) \rangle \varepsilon^*_\mu \bar{q}_\nu = 2f_{D^+\rho^+} (\varepsilon^* \cdot \bar{q} \varepsilon^* \cdot k - \varepsilon^* \cdot \bar{\varepsilon} \cdot k \cdot \bar{q}),$$

where $k + \bar{q}$ is the $D^+$ momentum, $k$ and $\varepsilon$ stand for the $\rho^+$ momentum and polarization vector, respectively, and the same constant $f_{D^+\rho^+}$ parametrizes form-factor effects at $q^2 = 0$ in the two equations, which are related [34] by virtue of the identity $2i \sigma^{\mu\nu} \gamma_5 = \varepsilon^{\mu\nu\rho\tau} \sigma_{\rho\tau}$. They lead to

$$\mathcal{M}_{D^+\rightarrow\rho^+\bar{\gamma}} = 4f_{D^+\rho^+} \varepsilon^* \bar{q}_\nu C + i (\varepsilon^* \cdot \bar{q} \varepsilon^* \cdot k - \varepsilon^* \cdot \bar{\varepsilon} \cdot k \cdot \bar{q}) C_5,$$

which is U(1)$_D$-gauge invariant and from which we obtain the branching fraction

$$\mathcal{B}(D^+ \rightarrow \rho^+\bar{\gamma}) = \frac{\tau_{D^+} f^2_{D^+\rho^+} (m^2_{D^+} - m^2_{\rho^+})^3}{2 \pi m^5_{D^+}}(|C|^2 + |C_5|^2),$$

where $\tau_{D^+}$ represents the lifetime of $D^+$ and $m_X$ denotes the mass of $X$. The branching fractions of $D^0 \rightarrow \rho^0\bar{\gamma}$, $\omega\bar{\gamma}$ and $D_s^+ \rightarrow K^+\bar{\gamma}$ have analogous expressions.

For numerical work, we employ

$$f_{D^+\rho^+} = \sqrt{2} f_{D^0\rho^0} = 0.658^{+0.038}_{-0.031}, \quad f_{D^0\omega} = 0.610^{+0.036}_{-0.030}, \quad f_{D_{K^+}^+K^{++}} = 0.639^{+0.042}_{-0.044},$$

(6)
which have been estimated in Ref. [33] with light-cone sum rules in the framework of heavy-quark effective field theory. The relation between $f_{D^+\rho^+}$ and $f_{D^0\rho}$ in Eq. (6) follows from the quark flavor contents $\rho^+ \sim u\bar{d}$ and $\rho^0 \sim (u\bar{u} - d\bar{d})/\sqrt{2}$. For comparison, an earlier analysis within a constituent quark model [36] yielded $f_{D^+\rho^+} = 0.66$ and $f_{D^+K^{*+}} = 0.71$ with uncertainties of around 10%, which are compatible with their newer counterparts in Eq. (6). Additional input parameters are the measured $D^{+\ast0}$ and $D_s^+$ lifetimes and masses and the light meson masses, namely $m_{\rho^+} = 775.11(34)$, $m_{\rho^0} = 775.26(25)$, $m_\omega = 782.65(12)$, and $m_{K^{*+}} = 895.5(8)$, in units of MeV, from Ref. [37]. Thus, the only remaining unknowns are the coefficients $C$ and $C_5$, which depend on the details of the NP model.

These processes can serve as a valuable tool in the search for the massless dark photon if their rates are not highly suppressed. Clearly, this can happen provided that one or both of $C$ and $C_5$ are not too small in size. It turns out that there is at least one NP model in the recent literature [7] which offers some nonnegligible viable values of these coefficients. Those numbers can then serve as benchmarks to illustrate how these charmed-hadron decays may probe the dark photon’s existence.

In the NP scenario of Ref. [7] the $c \to u\gamma$ operators arise from loop diagrams involving new particles consisting of massive fermions which are singlets under the SM gauge group and heavy new scalar bosons which carry some of the SM gauge charges. The new fermions and bosons are all charged under $U(1)_D$ and have Yukawa-like interactions with the $u$ and $c$ quarks. This allows for the construction of the dimension-five operators in Eq. (2). Since in this paper we are mainly interested in the implications of these effective interactions for the charmed-hadron decays, we will not dwell further on the specifics of the underlying NP model. Rather, we will simply adopt the relevant results available from Ref. [7] at face value and apply them to the determination of the rates of the charmed-hadron processes.

Particularly, as elaborated in Ref. [7], the branching fraction of the inclusive transition $c \to u\gamma$ could reach $\mathcal{B}(c \to u\gamma) \sim 10^{-4}$, with the pertinent constraints, from dark matter and vacuum saturation considerations, having been taken into account. Accordingly, for the purposes of our study we can take $\mathcal{B}(c \to u\gamma) < 5 \times 10^{-5}$, which translates into

$$|C|^2 + |C_5|^2 < \frac{2.0 \times 10^{-15}}{\text{GeV}^2}. \quad (7)$$

The rates of all the decays which will be discussed shortly are proportional to $|C|^2 + |C_5|^2$. Consequently, if future searches for these transitions come up with limits stronger than the branching fractions we predict below based on Eq. (7), the implied constraints apply only to this combination of the coefficients. To probe $C$ and $C_5$ separately would require measuring more complicated

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1 This has been found with $\mathcal{B}(c \to u\gamma) = 12 \mathcal{B}(c \to \ell^+X)\exp(1/|\Lambda_L|^2 + 1/|\Lambda_R|^2)/[G_F^2m_c^2|V_{cs}|^2f_1(m_\omega^2/m_c^2)]$ from [7] containing the usual Fermi constant $G_F$ and the inputs $\mathcal{B}(c \to \ell^+X)\exp = 0.096$, $m_c = 1.67$ GeV, $|V_{cs}| = 0.986$, and $f_1(m_\omega^2/m_c^2) \sim 1$ also provided therein, along with the relation $1/|\Lambda_L|^2 + 1/|\Lambda_R|^2 = 8|C|^2 + 8|C_5|^2$. 

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5
observables, such as the angular distributions of the decay products, but this step will be worth
taking only after any of these decays are discovered.

Incorporating Eq. (7) and the central values of the aforementioned input parameters into the
meson branching fractions, we arrive at

\[ \mathcal{B}(D^+ \to \rho^+ \gamma) < 8.1 \times 10^{-4}, \]
\[ \mathcal{B}(D^0 \to \rho^0 \gamma) < 1.6 \times 10^{-4}, \]
\[ \mathcal{B}(D^0 \to \omega \gamma) < 2.7 \times 10^{-4}, \]
\[ \mathcal{B}(D_s^+ \to K^{*+} \gamma) < 3.8 \times 10^{-4}. \]

Interestingly, these numbers exceed the current limit \( \mathcal{B}(D^0 \to \text{invisibles}) < 9.4 \times 10^{-5} \) at 90%
confidence level \[37\] set by the Belle Collaboration \[38\]. This suggests that one or more of
the results in Eq. (8) may be within the reaches of ongoing experiments such as Belle II \[30\] and
BESIII \[32\] and future facilities such as super charm-tau factories.

For \( D^0 \to \gamma \bar{\gamma} \), we need the matrix elements

\[ \langle \gamma(k)| \bar{u} \gamma_{\mu} c | D^0(k + \bar{q}) \rangle \bar{\varepsilon}_\mu \bar{q}_\nu = i e f_{D^0 \gamma} \epsilon^{\nu \rho \mu \nu} \bar{\varepsilon}_\rho \varepsilon^* \bar{q}_\nu, \]
\[ \langle \gamma(k)| \bar{u} \gamma_{5 \mu} \gamma_5 c | D^0(k + \bar{q}) \rangle \bar{\varepsilon}_\mu \bar{q}_\nu = e f_{D^0 \gamma}(\bar{\varepsilon}^* \cdot \bar{q} \varepsilon^* \cdot k - \bar{\varepsilon}^* \cdot \varepsilon^* \cdot k \bar{q}), \]

which have also been considered in the literature \[33, 40\] and involve the electric charge \( e \), the
photon’s polarization vector \( \bar{\varepsilon} \), and the form-factor parameter \( f_{D^0 \gamma} \). Therefore, the decay amplitude
satisfies both the U(1)\(_D\) and electromagnetic gauge symmetries. The corresponding branching
fraction is

\[ \mathcal{B}(D^0 \to \gamma \bar{\gamma}) = \frac{\alpha_e}{2} \tau_{D^0} \frac{f_{D^0 \gamma}^2 m_{D^0}^3}{m_{D^0}^3 (|C|^2 + |C_5|^2)}, \]

where \( \alpha_e = e^2/(4\pi) = 1/137 \). Including \( f_{D^0 \gamma} = 0.33 \) from Ref. \[40\] and Eq. (7) then leads to

\[ \mathcal{B}(D^0 \to \gamma \bar{\gamma}) < 3.2 \times 10^{-6}. \]

This is roughly two orders of magnitude below its counterparts in Eq. (8), which is due to the
presence of \( \alpha_e \).

III. DECAY OF BOTTOM, CHARMED MESON

The \( c \to u \bar{\gamma} \) transition can also bring about the decay of the bottom, charmed pseudoscalar-
meson \( B_c^+ \) into the bottom charged vector-meson \( B^{*+} \) plus the dark photon. Consequently, the
amplitude for \( B_c^+ \to B^{*+} \bar{\gamma} \) and its branching fraction are similar to the corresponding quantities
in the \( D^+ \to \rho^+ \bar{\gamma} \) case, given in Eqs. (4) and (5). Especially, based on the latter we have

\[ \mathcal{B}(B_c^+ \to B^{*+} \bar{\gamma}) = \frac{\tau_{B_c^+} f_{B_c^+ B^{*+}}^2 (m_{B_c^+}^2 - m_{B^{*+}}^2)^3}{2\pi m_{B_c^+}^3 (|C|^2 + |C_5|^2)} \]

(12)
To evaluate this, we use the $B_c^+$ lifetime and mass and $m_{B^+} = 5324.70(22)$ MeV from Ref. [37], as well as $f_{B_c B_s^0} = 0.23 \pm 0.04$ which has been computed in Ref. [41] with QCD sum rules and is in agreement with the older result $f_{B_c B_s^0} = 0.24$ of Ref. [42] from an application of the constituent quark model of Ref. [43]. Combining their central values and Eq. (7) with Eq. (12), we get

$$\mathcal{B}(B_c^+ \to B^{*+}\gamma) < 7.1 \times 10^{-5}.$$  \hspace{1cm} (13)

### IV. DECAYS OF SINGLY CHARMED BARYONS

In the baryon sector, the occurrence of the $c \to u\gamma$ transition in Eq. (2) will cause charmed baryons to undergo $\Delta C = 1$ decays into a lighter baryon and the dark photon. In this section we focus on the singly charmed baryons $\Lambda_c^+, \Xi_c^+$, and $\Xi^0_c$, which have spin parity $J^P = 1/2^+$, make up a flavor SU(3) antitriplet, and decay weakly [37]. We examine in particular the modes $\Lambda_c^+ \to p\gamma$ and $\Xi_c^{+0} \to \Sigma_c^{++}\gamma$.

To derive the amplitude $\mathcal{M}_{\Lambda_c^+ \to p\gamma}$ for $\Lambda_c^+ \to p\gamma$ from the short-distance interaction described by $\mathcal{L}_{uc\gamma}$, we need to know the matrix elements $\langle p|\bar{u}\sigma^{\mu\nu}c|\Lambda_c^+\rangle$ and $\langle p|\bar{u}\sigma^{\mu\nu}\gamma_5 c|\Lambda_c^+\rangle$. Contracting them with dark photon’s momentum $\vec{q}$ and polarization vector $\vec{\varepsilon}$, as explained in the appendix, we can write

$$\langle p|\bar{u}\sigma^{\mu\nu}(1, \gamma_5)\gamma^\nu\bar{c}\gamma_5 u|\Lambda_c^+\rangle \varepsilon^\mu q^\nu = f_{\Lambda_c^+ p} \bar{U}_p \sigma^{\mu\nu}(1, \gamma_5)U_{\Lambda_c^+} \varepsilon^\mu q^\nu,$$  \hspace{1cm} (14)

where $U_{\Lambda_c^+}$ and $U_p$ stand for the Dirac spinors of the baryons and the same constant $f_{\Lambda_c^+ p}$, which encodes form-factor effects at $q^2 = 0$, enters both of the matrix elements. It follows that

$$\mathcal{M}_{\Lambda_c^+ \to p\gamma} = 2 f_{\Lambda_c^+ p} \bar{U}_p (\bar{C} + \gamma_5 \bar{C}_5) i\sigma^{\mu\nu}U_{\Lambda_c^+} \varepsilon^\mu q^\nu,$$  \hspace{1cm} (15)

which fulfills U(1)$_D$-gauge invariance, and so the branching fraction is

$$\mathcal{B}(\Lambda_c^+ \to p\gamma) = \frac{\tau_{\Lambda_c^+} f_{\Lambda_c^+ p}^3 m_{\Lambda_c^+}^2 - m_p^2)^3}{2\pi m_{\Lambda_c^+}^3 \left(|C|^2 + |C_5|^2\right)},$$  \hspace{1cm} (16)

where $\tau_{\Lambda_c^+}$ is the $\Lambda_c^+$ lifetime. Numerically, we adopt $f_{\Lambda_c^+ p} = 0.50(1 \pm 0.07)$ from the lattice QCD calculation in Ref. [44]. Incorporating its central value and those of $\tau_{\Lambda_c^+}$, $m_{\Lambda_c^+}$, and $m_p$ from Ref. [37] into Eq. (16), we then translate the limit in Eq. (7) into

$$\mathcal{B}(\Lambda_c^+ \to p\gamma) < 1.7 \times 10^{-4}.$$  \hspace{1cm} (17)

If instead $f_{\Lambda_c^+ p} = 0.38(1 \pm 0.1)$, found in the relativistic quark model [45], we would reach a lower result, $\mathcal{B}(\Lambda_c^+ \to p\gamma) < 9.6 \times 10^{-5}$. These different numbers indicate the degree of uncertainty in the prediction.

In the case of $\Xi_c^0 \to \Sigma^0\gamma$, we employ the form-factor parameter $f_{\Xi_0^0 \Sigma^0} = 0.46 \pm 0.12$, which has been estimated in the framework of light-cone QCD sum rules [46]. Assuming isospin symmetry, we
also have $f_{\Xi^+ \Sigma^+} = f_{\Xi^0 \Sigma^0}$ for $\Xi_c^+ \to \Sigma^+ \gamma$. Their branching fractions have expressions obtainable from Eq. (16), with suitable modifications. Then, with Eq. (7) and the central values of $f_{\Xi^0 \Sigma^0}$ and the measured $\Xi_c^{+,0}$ lifetimes and masses and $\Sigma^{+,0}$ masses [37], we arrive at

$$
B(\Xi_c^+ \to \Sigma^+ \gamma) < 3.1 \times 10^{-4},
$$
$$
B(\Xi_c^0 \to \Sigma^0 \gamma) < 7.8 \times 10^{-5}.
$$

The difference between them is mainly due to $\tau_{\Xi_c^+} = 3.9 \tau_{\Xi_c^0}$.

V. DECAYS OF DOUBLY CHARMED BARYON

Only one doubly charmed baryon, the $\Xi_{cc}^{++}$, has been discovered so far [47], with its lifetime and mass now also known [37]. We take its spin parity, which is not yet determined experimentally, to be $J^P = 1/2^+$ based on quark-model expectations [47]. Here we consider the decay channels $\Xi_{cc}^{++} \to \Sigma_c(2455)^{++} \gamma$ and $\Xi_{cc}^{++} \to \Sigma_c(2520)^{++} \gamma$ arising from $\mathcal{L}_{cc\gamma}$, as some information on the pertinent baryonic form factors has recently become available [48].

Since $\Sigma_c(2520)^{++}$ also has $J^P = 1/2^+$ [37], the branching fraction of $\Xi_{cc}^{++} \to \Sigma_c(2455)^{++} \gamma$ is analogous to that in Eq. (16), namely

$$
B(\Xi_{cc}^{++} \to \Sigma_c(2455)^{++} \gamma) = \frac{\tau_{\Xi_{cc}^{++}} f_{\Xi_{cc}^{++} \Sigma_c(2455)}^2}{2\pi m_{\Xi_{cc}^{++}}^3} \left( \frac{m_{\Xi_{cc}^{++}}^2 - m_{\Sigma_c(2455)^{++}}^2}{m_{\Sigma_c(2455)^{++}}} \right)^3 (|C|^2 + |C_5|^2).
$$

where $m_{\Xi_{cc}^{++}} \equiv m_{\Xi_{cc}^{++}}$. In this equation, we use the form-factor parameter $f_{\Xi_{cc}^{++} \Sigma_c(2455)} = -0.798$ which has been computed in the light-front quark model with an uncertainty of several percent [48]. With the central values of the measured $\Xi_{cc}^{++}$ lifetime and mass and $m_{\Sigma_c(2455)^{++}} = 2453.97$ MeV from Ref. [37], plus the bound in Eq. (7), we then get

$$
B(\Xi_{cc}^{++} \to \Sigma_c(2455)^{++} \gamma) < 5.9 \times 10^{-4}.
$$

In the second channel, the daughter baryon $\Sigma_c(2520)^{++}$ has spin parity $J^P = 3/2^+$ [37]. As discussed in the appendix, the baryonic matrix elements in this case are

$$
\langle \Sigma_c(2520)^{++} | T^{\mu\nu} (1, \gamma_5) c | \Xi_{cc}^{++} \rangle \tilde{\varepsilon}_\mu^{\ast} \tilde{q}_\nu = \overline{U}_{\Sigma_c} (\gamma_5, 1) \left[ (\tilde{\tilde{F}} \tilde{q}_\kappa - \tilde{\tilde{\varepsilon}}_\kappa^{\ast} \tilde{q}_\kappa) \mathcal{F} + \bar{q}_\kappa \tilde{\tilde{F}} \tilde{\tilde{q}}_\kappa \tilde{F} \right] u_{\Xi_{cc}^{++}},
$$

where $U_{\Sigma_c}$ denotes the Rarita-Schwinger spinor of $\Sigma_c(2520)^{++}$ and the constants $\mathcal{F}$ and $\tilde{F}$ parametrize form-factor effects at $\bar{q}_\kappa^2 = 0$. We can then write the decay amplitude

$$
\mathcal{M}_{\Xi_{cc}^{++} \to \Sigma_c(2520)^{++} \gamma} = 2 \overline{U}_{\Sigma_c} (\gamma_5 C + C_5) \left[ (\tilde{\tilde{F}} \tilde{q}_\kappa - \tilde{\tilde{\varepsilon}}_\kappa^{\ast} \tilde{q}_\kappa) \mathcal{F} + \bar{q}_\kappa \tilde{\tilde{F}} \tilde{\tilde{q}}_\kappa \tilde{F} \right] u_{\Xi_{cc}^{++}},
$$

\[8\]
which is U(1)$_D$-gauge invariant and leads to the branching fraction

$$
\mathcal{B}(\Xi^{++}_{cc} \rightarrow \Sigma_c(2520)^{++}\bar{\gamma}) = \frac{\mathcal{E}_{\Xi^{++}_{cc}} \Delta \mathcal{F}^2}{12\pi m^2_{\Xi_{cc}} m^2_\gamma} \left(1 + \frac{4m^2_\Sigma}{\Delta^2} + \frac{2\tilde{\mathcal{F}}}{\mathcal{F}} + \frac{\Delta^2\tilde{\mathcal{F}}^2}{m^2_{\Xi_{cc}} \mathcal{F}^2}\right)\left(|\mathcal{C}|^2 + |\mathcal{C}_5|^2\right),
$$

(23)

where $m_{\Sigma_c} \equiv m_{\Sigma_c(2520)^{++}}$ and $\Delta = \left(m^2_{\Xi_{cc}} - m^2_{\Sigma_c}\right)^{1/2}$. Adopting $\mathcal{F} = -0.635$ and $\tilde{\mathcal{F}} = 0.330$, which have been estimated in the light-front quark model with uncertainties of several percent[48], with $m_{\Sigma_c} = 2518.41$ MeV[37] and Eq. (7), we then find

$$
\mathcal{B}(\Xi^{++}_{cc} \rightarrow \Sigma_c(2520)^{++}\bar{\gamma}) < 2.2 \times 10^{-4}.
$$

(24)

VI. CONCLUSIONS

If a dark sector beyond the SM exists, it may contain a new unbroken Abelian gauge symmetry under which all SM fields are singlets. Furthermore, its associated gauge boson, the massless dark photon, could have nonnegligible flavor-changing dipole-type couplings to the $u$ and $c$ quarks. If this is realized in nature, it will bring about the FCNC decays of charmed hadrons into a lighter hadron plus missing energy carried away by the dark photon. Entertaining these possibilities, we have investigated a number of such $\Delta C = 1$ reactions. We study particularly those in which the parent particles are the charmed pseudoscalar-mesons $D^{+,0}$ and $D^+_s$, the bottom, charmed pseudoscalar-meson $B^+_c$, the singly charmed baryons $\Lambda^+_c$, $\Xi^+_c$, and $\Xi^0_c$, and the doubly-charmed baryon $\Xi^{++}_{cc}$. We also look at $D^0 \rightarrow \gamma\bar{\gamma}$ with the ordinary photon besides $\bar{\gamma}$ in the final state. In the context of a simplified new-physics model, we show that the majority of their branching fractions are permitted by present constraints to be as high as a few times $10^{-4}$. Such numbers may be within the sensitivity reaches of the ongoing Belle II and BESIII experiments. Since the same one or two $c \rightarrow u\bar{\gamma}$ operators are responsible for all these decays, detecting one of them automatically implies specific predictions for the others, allowing for additional empirical tests on the dark photon scenario. Conversely, measuring a bound on one of the decays will translate into expected bounds on the others. The results of this work will hopefully help motivate dedicated attempts to hunt massless dark photons via charmed-hadron processes. These efforts will be complementary to the quests for the massive dark photon, which is phenomenologically very distinct from the massless one.

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2 It is worth noting that Eqs. (22) and (23) can be shown to be consistent in form with the amplitude and branching fraction, respectively, of $\Lambda_b \rightarrow \Lambda(1520)\gamma$, which involves the ordinary photon, in[49], after interchanging the parity-conserving and -violating terms in the amplitude, as $\Sigma_c(2520)^{++}$ and $\Lambda(1520)$ are opposite in parity.
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Appendix A: Baryonic matrix elements of $\bar{u}\sigma^{\mu\nu}(1,\gamma_5)c$

For the $\Delta C = 1$ transition between baryons $\mathcal{B}$ and $\mathcal{B}'$ with spin parity $J^P = 1/2^+$, the baryonic matrix element of $\bar{u}\sigma^{\mu\nu}c$ has the general expression \cite{50,51}

$$\langle \mathcal{B}'(k) | \bar{u}\sigma^{\mu\nu}c | \mathcal{B}(k + \bar{q}) \rangle = \bar{U}' \left( \sigma^{\mu\nu} f_1 - \gamma^{[\mu} k^{\nu]} i f_2 - \gamma^{[\mu} q^{\nu]} i f_3 - k^{[\mu} q^{\nu]} i f_4 \right) U,$$  \hspace{1cm} (A1)

where $k + \bar{q}$ and $k$ are the momenta of $\mathcal{B}$ and $\mathcal{B}'$, respectively, $U$ and $U'$ represent their Dirac spinors, $f_{1,2,3,4}$ denote form factors which are functions of $\bar{q}^2$, and $X^{[\mu\nu]} \equiv X^{\mu\nu} - X^{\nu\mu}$. With the aid of the identity $2i\sigma_{\nu\tau} = \epsilon_{\nu\mu\tau\sigma}$ for $\epsilon_{0123} = +1$, we can obtain from Eq. (A1) the corresponding matrix element of $\bar{u}\sigma_{\nu\tau}\gamma_5 c$,

$$\langle \mathcal{B}'(k) | \bar{u}\sigma_{\nu\tau}\gamma_5 c | \mathcal{B}(k + \bar{q}) \rangle = \bar{U}' \left[ \sigma_{\nu\tau} \gamma_5 f_1 - \gamma_{\nu\tau} \left( \frac{1}{2} \gamma^{[\mu} k^{\nu]} f_2 + \gamma^{[\mu} q^{\nu]} f_3 + k^{[\mu} q^{\nu]} f_4 \right) \right] U. \hspace{1cm} (A2)$$

For application in the amplitude for $\mathcal{B} \rightarrow \mathcal{B}'\bar{\gamma}$, the matrix elements in Eqs. (A1) and (A2) are to be contracted with the dark photon’s polarization vector $\varepsilon$ and momentum $\bar{q}$. Thus, imposing $\varepsilon \cdot \bar{q} = 0$ and $\bar{q}^2 = \bar{q}\bar{q} = 0$, after straightforward algebra we arrive at

$$\langle \mathcal{B}'(k) | \bar{u}\sigma^{\mu\nu}c | \mathcal{B}(k + \bar{q}) \rangle \varepsilon^*_\mu \bar{q}_\nu = \bar{U}' \left\{ \sigma^{\mu\nu} f_1^{(0)} - \frac{i}{2} \left[ \gamma^{[\mu} \gamma^{\nu]} (k + \bar{q}) - k^{[\mu} \gamma^{,\nu]} \right] f_2^{(0)} \right\} U \varepsilon^*_\mu \bar{q}_\nu,$$

$$= f_{\mathcal{B}\mathcal{B}'} \bar{U}' \sigma^{\mu\nu} U \varepsilon^*_\mu \bar{q}_\nu, \hspace{1cm} (A3)$$

$$\langle \mathcal{B}'(k) | \bar{u}\sigma^{\mu\nu}\gamma_5 c | \mathcal{B}(k + \bar{q}) \rangle \varepsilon^*_\mu \bar{q}_\nu = \bar{U}' \left\{ \sigma^{\mu\nu} \gamma_5 f_1^{(0)} - \frac{i}{2} \epsilon^{\mu\nu\tau\sigma} \left[ \gamma_5 \gamma_\tau (k + \bar{q}) - k^{[\mu} \gamma_\sigma \right] f_2^{(0)} \right\} U \varepsilon^*_\mu \bar{q}_\nu,$$

$$= f_{\mathcal{B}\mathcal{B}'} \bar{U}' \sigma^{\mu\nu} \gamma_5 U \varepsilon^*_\mu \bar{q}_\nu, \hspace{1cm} (A4)$$

where $f_{1,2}^{(0)}$ stand for $f_{1,2}$ evaluated at $\bar{q}^2 = 0$ and

$$f_{\mathcal{B}\mathcal{B}'} = f_1^{(0)} + f_2^{(0)} \left( m_{\mathcal{B}'} - m_{\mathcal{B}} \right), \hspace{1cm} (A5)$$

with $m_{\mathcal{B}(c)}$ being the mass of $\mathcal{B}(c)$. These results lead to Eq. (A4). Evidently, the same combination, $f_{\mathcal{B}\mathcal{B}'}$, of form factors at $\bar{q}^2 = 0$ enters both Eqs. (A3) and (A4), in agreement with what was previously found \cite{52,53} concerning the baryonic matrix elements of these types of tensor operators.

In the case where $\mathcal{B}'$ is replaced by a spin-3/2 baryon $\mathcal{B}'$, we can write

$$\langle \mathcal{B}'(k) | \bar{u}\sigma^{\mu\nu}c | \mathcal{B}(k + \bar{q}) \rangle = \left\{ \bar{U}' \left[ \epsilon^{\mu\nu} g_1 + \bar{U}' \gamma_5 g_2 + \bar{U}' \gamma_\eta \left( i \sigma^{\mu\nu} g_3 + \gamma^{[\mu} k^{\nu]} g_4 \right) + \bar{U}' \gamma^{[\mu} q^{\nu]} g_5 + \bar{U}' \gamma_\eta \left( \gamma^{[\mu} q^{\nu]} g_6 + k^{[\mu} q^{\nu]} g_7 \right) \right\} \gamma_5 U,$$  \hspace{1cm} (A6)
where $k$ and $\mathbb{U}^\mu$ are, respectively, the momentum of $B'$ and its Rarita-Schwinger spinor, the latter satisfying the requirements $\gamma_\mu \mathbb{U}^\mu = k_\mu \mathbb{U}^\mu = 0$, and $g_{1,2,\ldots,7}$ represent form factors which depend on $q^2$. It follows that

$$
\langle B'(k)|\bar{u}\sigma^{\mu\nu}c|B(k+\hat{q})\rangle \varepsilon^*_\mu \bar{q}_\nu = \mathbb{U}^\mu \gamma^\nu \varepsilon^*_\mu \bar{q}_\nu \left[g_1^{(0)} + \frac{1}{2} (m_{B'} - m_B) g_2^{(0)} \right] \gamma_5 \mathbb{U}
+ \mathbb{U}^\mu \bar{q}_\eta \not\!\eta \left[\varepsilon^*_\mu \bar{q}_\nu - g_3^{(0)} - \frac{1}{2} (m_{B'} + m_B) g_4^{(0)} \right] \gamma_5 \mathbb{U}
= \mathbb{U}^\mu \left[ (\varepsilon^*_\mu \not\!\eta - \not\!\bar{q}_\eta \gamma_5 \not\!\nu) \mathcal{F} + \frac{\bar{q}_\eta \not\!\gamma_5 \not\!\nu \mathcal{F}}{m_{B'}} \right] \gamma_5 \mathbb{U},
$$

(A7)

where $g_j^{(0)}$ denotes the value of $g_j$ at $q^2 = 0$ and

$$
\mathcal{F} = g_1^{(0)} + \frac{g_2^{(0)}}{2} (m_{B'} - m_B),
\frac{\mathcal{F}}{m_{B'}} = g_2^{(0)} \frac{g_3^{(0)}}{2} - g_3^{(0)} - \frac{g_4^{(0)}}{2} (m_{B'} + m_B),
$$

(A8)

with $m_{B'}$ being the mass of $B'$. Moreover,

$$
\langle B'(k)|\bar{u}\sigma^{\mu\nu}\gamma_5 c|B(k+\hat{q})\rangle \varepsilon^*_\mu \bar{q}_\nu = \frac{\epsilon^{\mu\nu\kappa\tau}}{2} \langle B(k)|\bar{u}\sigma_{\kappa\tau} c|B(k+\hat{q})\rangle \varepsilon^*_\mu \bar{q}_\nu
= \mathbb{U}^\mu \left[ (\varepsilon^*_\mu \not\!\eta - \not\!\bar{q}_\eta \gamma_5 \not\!\nu) \mathcal{F} + \frac{\bar{q}_\eta \not\!\gamma_5 \not\!\nu \mathcal{F}}{m_{B'}} \right] U.
$$

(A9)

From Eqs. (A7) and (A9), we get Eq. (21).

In the derivation of the $B \rightarrow B'\gamma$ rate, the absolute square of the decay amplitude needs to be summed over the initial and final particles’ polarizations. The expression for the sum over the $B'$ polarizations is available from the literature (e.g. [54]). It is given here for completeness:

$$
\sum_{\gamma = -\frac{3}{2}}^{\gamma = \frac{3}{2}} \mathcal{U}_B^\gamma (k; \zeta) \mathcal{U}_B^\gamma (k; \zeta) = (k + m_{B'}) \left( \frac{\gamma_\mu \gamma_\nu}{3} \mathcal{G}^{\mu\nu}_B (k) \mathcal{G}^{\gamma\eta}_B (k) - \mathcal{G}^{\gamma\eta}_B (k) \right),
$$

(A10)

where $\mathcal{G}^{\mu\nu}_B (k) = g^{\mu\nu} - k_\mu k_\nu/m_{B'}^2$.

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