Charge asymmetry in $K^{\pm}\to\pi^{\pm}\gamma\gamma$ induced by the electromagnetic penguin operators

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The CP-violating charge asymmetry in the decays $K^{\pm}\to\pi^{\pm}\gamma\gamma$, which is induced by the electromagnetic penguin operators, has been studied both in the standard model and in its extensions. Because of a large enhancement of the Wilson coefficients of the electromagnetic penguin operators in the supersymmetric extensions of the standard model, a significant upper bound on this charge asymmetry could be expected, and thus high-precision measurements of this interesting CP-violating quantity might explore new physics effects beyond the standard model.

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Rare kaon decays provide a very useful laboratory both to test the standard model (SM) and to explore new physics beyond it [1–4]. It is of particular interest to study the CP-violating effects, which arise in such weak decays from dimension-five operators, including the electromagnetic and chromomagnetic penguin operators (EMO and CMO) since the CP violation induced by these operators is suppressed in the SM; however, it could be enhanced in its extensions [5–12]. On the other hand, present experiments, HyperCP [13], and future NA48 experiments [14], are going to substantially improve the present limits on the Wilson coefficients of these operators by studying CP-violating charge asymmetries in charged kaon decays, such as $K^{\pm}\to(3\pi)^{\pm}$, $K^{\pm}\to\pi^{\pm}\ell^{+}\ell^{-}$, as well as one-photon or two-photon radiative decays. It is expected that charged kaon decays could be an ideal framework to explore direct CP violation, or CP violation of pure $\Delta S=1$ origin [3,6,7,11,12,15]. The purpose of this paper is devoted to the analysis of the CP-violating charge asymmetry induced by the EMO in the two-photon radiative charged kaon decay, $K^{\pm}\to\pi^{\pm}\gamma\gamma$, both in the SM and in its possible extensions.

The general weak effective Hamiltonian, contributed by the EMO and CMO, can be written as [5]

$$\mathcal{H}_{\rm eff}=C_{\gamma}^{+}(\mu)Q_{\gamma}^{+}(\mu)+C_{\gamma}^{-}(\mu)Q_{\gamma}^{-}(\mu)+C_{g}^{+}(\mu)Q_{g}^{+}(\mu)+C_{g}^{-}(\mu)Q_{g}^{-}(\mu)+\text{H.c.},\tag{1}$$

where $C_{\gamma,g}^{\pm}$ are the Wilson coefficients and

$$Q_{\gamma}^{\pm}=\frac{eQ_d}{16\pi^2}(\overline{s_L}\sigma_{\mu\nu}d_R\pm\overline{s_R}\sigma_{\mu\nu}d_L)F^{\mu\nu},\tag{2}$$

$$Q_{g}^{\pm}=\frac{g_s}{16\pi^2}(\overline{s_L}\sigma_{\mu\nu}t_a d_R\pm\overline{s_R}\sigma_{\mu\nu}t_a d_L)G_a^{\mu\nu}.\tag{3}$$

Here $Q_d=-1/3$, and $\sigma_{\mu\nu}=i/2[\gamma_\mu,\gamma_\nu]$. Note that Eq. (1) with complex Wilson coefficients $C_{\gamma,g}^{\pm}$ could lead to new flavor structures beyond the SM, which generally depart from minimal flavor violation [16]. It is easy to see that the SM structure, $SU(2)_L\times U(1)_Y$, will impose the following chiral suppression for these operators [17,18]:

$$\mathcal{H}_{\rm eff}^{\rm SM}=\frac{G_F}{\sqrt{2}}V_{td}V_{ts}^{*}\left[C_{11}\frac{g_s}{8\pi^2}(m_d\overline{s_L}\sigma_{\mu\nu}t_a d_R+m_s\overline{s_R}\sigma_{\mu\nu}t_a d_L)G_a^{\mu\nu}+C_{12}\frac{e}{8\pi^2}(m_d\overline{s_L}\sigma_{\mu\nu}d_R+m_s\overline{s_R}\sigma_{\mu\nu}d_L)F^{\mu\nu}\right]+\text{H.c.},\tag{4}$$

with

$$C_{11}(m_W)=\frac{3x^2}{2(1-x)^4}\ln x-\frac{x^3-5x^2-2x}{4(1-x)^3},\tag{5}$$

$$C_{12}(m_W)=\frac{x^2(2-3x)}{2(1-x)^4}\ln x-\frac{8x^3+5x^2-7x}{12(1-x)^3},\tag{6}$$

where $x=m_t^2/m_W^2$ and t_a are the $SU(3)$ matrices. However, as we shall see, new flavor structures, for instance, from the supersymmetric extensions of the SM, allow us to avoid the chiral suppression for the operators in Eq. (4) [19].

It has been known that the $K^{\pm}\to\pi^{\pm}\gamma\gamma$ transitions are dominated by long-distance effects [20,21]. Within the framework of chiral perturbation theory [20,22], $K^{\pm}\to\pi^{\pm}\gamma\gamma$ receives the first nonvanishing contribution at $O(p^4)$ including both loops and anomalous and nonanomalous couterterms. The general $O(p^4)$ amplitude of the decay can be decomposed in the following way [20]:

$$M[K^{+}(k)\to\pi^{+}(p)\gamma(q_1,\epsilon_1)\gamma(q_2,\epsilon_2)]=\epsilon_{1\mu}(q_1)\epsilon_{2\nu}(q_2)\left[\frac{A(y,z)}{m_K^2}(q_2^{\mu}q_1^{\nu}-q_1\cdot q_2 g^{\mu\nu})+\frac{C(y,z)}{m_K^2}\varepsilon^{\mu\nu\alpha\beta}q_{1\alpha}q_{2\beta}\right],\tag{7}$$

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with

$$y=\frac{k\cdot(q_1-q_2)}{m_K^2},\quad z=\frac{(q_1+q_2)^2}{m_K^2}. \tag{8}$$

The physical region in the dimensionless variables y and z is given by

$$0\leqslant|y|\leqslant\frac{1}{2}\lambda^{1/2}(1,z,r_\pi^2),\quad 0\leqslant z\leqslant(1-r_\pi)^2, \tag{9}$$

where $r_\pi=m_\pi/m_K$ and $\lambda(a,b,c)=a^2+b^2+c^2-2(ab+ac+bc)$. Note that the invariant amplitudes $A(y,z)$ from loops and nonanomalous counterterms, and $C(y,z)$ from anomalous counterterms have to be symmetric under the interchange of q_1 and q_2 as required by Bose symmetry. In the SM, the $O(p^4)$ amplitude for $A(y,z)$ has been given in Ref. [20], which is

$$A(y,z)=\frac{G_8m_K^2\alpha_{\rm EM}}{2\pi z}\left[(r_\pi^2-1-z)F\left(\frac{z}{r_\pi^2}\right)+(1-r_\pi^2-z)F(z)+\hat{c}z\right], \tag{10}$$

where $\alpha_{\rm EM}=e^2/4\pi$, and $|G_8|=9.2\times10^{-6}\ \mathrm{GeV}^{-2}$. We do not display the explicit expression for $C(y,z)$ since it is irrelevant to the present discussion. $F(z/r_\pi^2)$ and $F(z)$ are generated from π and K loop diagrams respectively, which could be defined as

$$F(x)=\begin{cases}1-\dfrac{4}{x}\arcsin^2\left(\dfrac{\sqrt{x}}{2}\right), & x\leqslant4,\\ 1+\dfrac{1}{x}\left(\ln\dfrac{1-\sqrt{1-4/x}}{1+\sqrt{1-4/x}}+i\pi\right)^2, & x\geqslant4.\end{cases} \tag{11}$$

$\hat{c}$ in Eq. (10) is from $O(p^4)$ nonanomalous local counterterms,

$$\hat{c}=\frac{128\pi^2}{3}[3(L_9+L_{10})+N_{14}-N_{15}-2N_{18}], \tag{12}$$

where L_9 and L_{10} are couplings in the $O(p^4)$ strong chiral Lagrangian [23] and N_i ($i=14$, 15, 18) are couplings in the $O(p^4)$ weak chiral Lagrangian [22].

From Eq. (11) it is obvious that the π-loop contribution, which is proportional to $F(z/r_\pi^2)$ in Eq. (10), will generate a CP invariant absorptive part. Thus if $\hat{c}$ has a nonvanishing phase, the interference between these two parts will lead to the charge asymmetry in $K^\pm\to\pi^\pm\gamma\gamma$ as follows [20,24]:

$$\frac{\delta\Gamma}{2\Gamma}=\frac{|\Gamma(K^+\to\pi^+\gamma\gamma)-\Gamma(K^-\to\pi^-\gamma\gamma)|}{\Gamma(K^+\to\pi^+\gamma\gamma)+\Gamma(K^-\to\pi^-\gamma\gamma)}, \tag{13}$$

$$\begin{aligned}\delta\Gamma&=|\Gamma(K^+\to\pi^+\gamma\gamma)-\Gamma(K^-\to\pi^-\gamma\gamma)|\\&=\frac{\mathrm{Im}\,\hat{c}|G_8\alpha_{\rm EM}|^2m_K^5}{2^{10}\pi^5}\int_{4r_\pi^2}^{(1-r_\pi)^2}dz\lambda^{1/2}(1,z,r_\pi^2)\\&\quad\times(r_\pi^2-1-z)z\,\mathrm{Im}\,F(z/r_\pi^2).\end{aligned} \tag{14}$$

Thus information about the imaginary part of $\hat{c}$ is relevant to the estimate of $\delta\Gamma$. It is clear that L_9 and L_{10} cannot contribute to it. In the SM, $\mathrm{Im}\,\hat{c}$ is therefore governed by the weak-coupling combination $N_{14}-N_{15}-2N_{18}$ due to the Cabbibo-Kobayashi-Maskawa (CKM) phase [27]. This point has been explored in the past literature, and some small values of the charge asymmetry in $K^\pm\to\pi^\pm\gamma\gamma$ have been obtained in Refs. [24,3] and [25], respectively;

$$\left(\frac{\delta\Gamma}{2\Gamma}\right)^{\rm SM}\leqslant10^{-3}\quad[24], \tag{15}$$

and

$$\left(\frac{\delta\Gamma}{2\Gamma}\right)^{\rm SM}<10^{-4}\quad[3]. \tag{16}$$

We would like to give some remarks here. First, Eq. (15) is the updated version of the result given in Ref. [20], and a vanishing imaginary part of $N_{14}-N_{15}-2N_{18}$ has been predicted by the authors of Ref. [28] using $1/N_C$ analysis. Second, in deriving Eq. (16), the $O(p^6)$ unitarity corrections [25,26] from the physical $K^\pm\to3\pi$ vertex to $K^\pm\to\pi^\pm\gamma\gamma$ have been taken into account. As an order-of-magnitude estimate for the charge asymmetry in $K^\pm\to\pi^\pm\gamma\gamma$ induced by the EMO, in the present paper we do not consider the $O(p^6)$ unitarity contributions since, as pointed out in Ref. [25], it does not alter significantly the value of the charge asymmetry obtained at $O(p^4)$ in chiral perturbation theory.

Let us now try to delve into the analysis of the CP-violating charge asymmetry in $K^\pm\to\pi^\pm\gamma\gamma$ arising from the EMO formulated in the general effective Hamiltonian of Eq. (1). It will be shown below that this asymmetry is very small in the SM, however, it could be significant in the supersymmetric scenarios beyond the SM by comparison with those given in Eqs. (15) and (16). Using the same way given in Ref. [9], we first construct the effective Lagrangian that represents the EMO. In fact there are many possible chiral realizations of these operators. To the purpose of this paper, the leading order $O(p^4)$ realizations of the EMO, which are relevant to $K\to\pi\gamma\gamma$ transitions, could be expressed as

$$\mathcal{L}=a_1\mathcal{L}_1+a_2\mathcal{L}_2, \tag{17}$$

$$\begin{aligned}\mathcal{L}_1=\frac{eQ_d}{16\pi^2}C_\gamma^\pm\langle\lambda U(F_{L\mu\nu}+U^\dagger F_{R\mu\nu}U)\pm\lambda\\ \times(F_{L\mu\nu}+U^\dagger F_{R\mu\nu}U)U^\dagger\rangle F^{\mu\nu}+\mathrm{H.c.},\end{aligned} \tag{18}$$

and

$$\mathcal{L}_2 = i\frac{eQ_d}{16\pi^2}C_\gamma^{\pm}\langle \lambda U L_\mu L_\nu \pm \lambda L_\mu L_\nu U^\dagger\rangle F^{\mu\nu} + \text{H.c.}, \quad (19)$$

where a_1 and a_2 are unknown coupling constants, $L_\mu = iU^\dagger D_\mu U$, and $(\lambda)_{ij} = \delta_{3i}\delta_{2i}$. We use the standard notation in chiral perturbation theory [29], $F_{L\mu\nu} = F_{R\mu\nu} = eQF_{\mu\nu}$, $D_\mu U = \partial_\mu U - ie[Q,U]A_\mu$, and $Q = \mathrm{diag}(2,-1,-1)/3$. U is a unitary 3×3 matrix with $\det U = 1$, which collects the Goldstone meson fields (π, K, and η) as follows:

$$U = \exp(i\sqrt{2}\Phi/f_\pi),$$

$$\Phi = \frac{1}{\sqrt{2}}\lambda^a\phi_a(x) = \begin{pmatrix} \frac{\pi^0}{\sqrt{2}} + \frac{\eta_8}{\sqrt{6}} & \pi^+ & K^+ \\ \pi^- & -\frac{\pi^0}{\sqrt{2}} + \frac{\eta_8}{\sqrt{6}} & K^0 \\ K^- & \bar{K}^0 & -\frac{2\eta_8}{\sqrt{6}} \end{pmatrix}, \quad (20)$$

where the λ^a's are 3×3 Gell-Mann matrices and $f_\pi \simeq 93$ MeV. Note that our result [Eq. (17)] is not contrary to that in Ref. [28] since it does not give any contribution to the imaginary part of the weak-coupling combination $N_{14} - N_{15} - 2N_{18}$. However, it is easy to see that the effective Lagrangian [Eq. (17)] will give new contributions to decays $K^+\to\pi^+\gamma\gamma$ and $K^-\to\pi^-\gamma\gamma$, respectively, one can therefore expect that the new structures in Eqs. (18) and (19) with the complex Wilson coefficients $C_\gamma^{\pm}$ will induce a possible charge asymmetry in $K^{\pm}\to\pi^{\pm}\gamma\gamma$ decays. Consequently, the interference between the amplitude for $K^{\pm}\to\pi^{\pm}\gamma\gamma$ from Eq. (17) and the absorptive part in Eq. (10) will give

$$\begin{aligned}\delta\Gamma^{\text{EMO}} &= |\Gamma(K^+\to\pi^+\gamma\gamma) - \Gamma(K^-\to\pi^-\gamma\gamma)|^{\text{EMO}} \\ &= \frac{\alpha_{\text{EM}}^2|G_8|m_K^5}{3\cdot 2^8\pi^5 f_\pi^2}\left|\text{Im}\, C_\gamma^+\left(\frac{2}{3}a_1 - a_2\right)\right| \\ &\quad \times \int_{4r_\pi^2}^{(1-r_\pi)^2} dz\,\lambda^{1/2}(1,z,r_\pi^2)(r_\pi^2 - 1 - z)z \\ &\quad \times \text{Im}\, F(z/r_\pi^2). \end{aligned} \quad (21)$$

The first observation of $K^+\to\pi^+\gamma\gamma$ has been reported by the BNL E787 Collaboration [21], and the branching ratio of the decay has been measured [21,30],

$$\text{Br}(K^+\to\pi^+\gamma\gamma) = (1.1\pm 0.3\pm 0.1)\times 10^{-6}. \quad (22)$$

Thus we have

$$\left(\frac{\delta\Gamma}{2\Gamma}\right)^{\text{EMO}} = (2.4\pm 0.7)\times 10^6\left|\text{Im}\ C_\gamma^+\left(\frac{2}{3}a_1 - a_2\right)\right|. \quad (23)$$

Our next task is to evaluate the magnitude of $\text{Im}\, C_\gamma^+$ and $(2/3a_1 - a_2)$ to check whether it is possible to get a significant charge asymmetry of $K^{\pm}\to\pi^{\pm}\gamma\gamma$ from Eq. (23). However, since unknown constants a_1 and a_2 in Eqs. (18) and (19) are related to the low-energy chiral dynamics, at the present we have no model-independent way to give a reliable determination of them. In the following we will estimate them using naive dimensional analysis, within the chiral quark model, and employing lattice calculation, respectively.

Naive dimensional analysis. As the order-of-magnitude estimate, using naive dimensional analysis [31], we can obtain

$$a_1 \sim a_2 \sim f_\pi\frac{f_\pi}{\Lambda_\chi}, \quad (24)$$

with $\Lambda_\chi = 4\pi f_\pi$ as the chiral symmetry spontaneously breaking scale. So we get

$$\left|\left(\frac{2}{3}a_1 - a_2\right)\right| \sim a_1 \sim a_2 \sim \frac{f_\pi}{4\pi} \quad (25)$$

and

$$\left(\frac{\delta\Gamma}{2\Gamma}\right)^{\text{EMO}} \sim 1.8\times 10^4\ \ \text{GeV}|\text{Im}\, C_\gamma^+|. \quad (26)$$

The chiral quark model. The chiral quark model has been extensively used to study low-energy hadronic physics involving strong and weak interactions [32–35]. In order to study the direct CP violation in decays $K\to\pi\ell^+\ell^-(\ell = e,\ \mu)$, this model has been employed by the authors of Ref. [12] to evaluate the Bosonization of the EMO, which corresponds to the a_2 part in Eq. (17). In the same way as that in Ref. [12], one can obtain the effective Lagrangian corresponding to the a_1 part in Eq. (17). This leads to

$$a_1 = \frac{f_\pi^2}{4M_Q}, \quad a_2 = \frac{3M_Q}{8\pi^2}, \quad (27)$$

where the constituent quark mass M_Q could be set about 0.3 GeV [12]. Thus we get

$$\left(\frac{\delta\Gamma}{2\Gamma}\right)^{\text{EMO}} = (1.6\pm 0.4)\times 10^4\ \ \text{GeV}|\text{Im}\, C_\gamma^+|. \quad (28)$$

Lattice calculation. So far there is no direct lattice calculation on a_1 and a_2. However, note that the a_2 part in Eq. (17) will also contribute to transition $K\to\pi\gamma^*\to\pi\ell^+\ell^-$, and the first lattice calculation of the matrix element of the EMO, $\langle\pi^0|Q_\gamma^+|K^0\rangle$, has been done in Ref. [36]. One therefore could determine the value of a_2 by comparing the result from Eq. (17) and that from the lattice calculation. In general, the matrix element of the EMO can be parametrized in terms of a suitable parameter B_T [5,36]:

$$\langle \pi^0 | Q_\gamma^+ | K^0 \rangle = i \frac{\sqrt{2} e Q_d}{16\pi^2 m_K} B_T p_\pi^\mu p_K^\nu F_{\mu\nu}. \tag{29}$$

On the other hand, using the effective Lagrangian in Eq. (17), one can get

$$\langle \pi^0 | Q_\gamma^+ | K^0 \rangle = i \frac{\sqrt{2} e Q_d}{16\pi^2 m_K} \frac{2 a_2 m_K}{f_\pi^2} p_\pi^\mu p_K^\nu F_{\mu\nu}. \tag{30}$$

Since $B_T = 1.18 \pm 0.09$ has been found in the lattice calculation [36], from Eqs. (29) and (30), we have

$$a_2 = 0.010 \pm 0.001 \ \ \text{GeV}. \tag{31}$$

Unfortunately, now we are not able to use the similar way to extract any information on a_1 from the lattice calculation. However, by comparing the value of a_2 in Eq. (31) with those from naive dimensional analysis [in Eq. (24)] and the chiral quark model [in Eq. (27)], one can find that they are of the same order of magnitude. Meanwhile, a_1 from Eqs. (24) and (27) are also of the same order of magnitude. This leads to the same order of magnitude estimates for the charge asymmetry in Eqs. (26) and (28), which are from naive dimensional analysis and the chiral quark model, respectively. Therefore, in general,

$$\left(\frac{\delta\Gamma}{2\Gamma}\right)^{\text{EMO}} \sim 1.0 \times 10^4 \ \ \text{GeV} |\text{Im}\, C_\gamma^+| \tag{32}$$

could be expected except in a fine-tuning case in which there is an accidental cancellation between $2/3a_1$ and a_2 (since we cannot reliably fix the relative sign of a_1 and a_2 in a model-independent way).

In order to go further into the charge asymmetry in Eq. (32), now one has to compute the imaginary parts of the Wilson coefficients $C_\gamma^\pm$, which are related to the short-distance physics. In the SM it is easy to get $\text{Im}\, C_\gamma^+$ from Eqs. (1) and (4) as

$$|\text{Im}\, C_\gamma^+|^{\text{SM}} = \frac{3 G_F}{\sqrt{2}} (m_s + m_d) |\text{Im}\, \lambda_t C_{12}|, \tag{33}$$

where $\lambda_t = V_{td} V_{ts}^*$. Due to the smallness of $\text{Im}\, \lambda_t \sim 10^{-4}$, this contribution from the SM to the charge asymmetry in $K^\pm \to \pi^\pm \gamma\gamma$ is strongly suppressed and could be negligible. Therefore in the following we will turn our attention to physics beyond the SM.

Among the possible new physics scenarios, low-energy supersymmetry (SUSY) [37] represents one of the most interesting and consistent extensions of the SM. In generic supersymmetric models, the large number of new particles carrying flavor quantum numbers would naturally lead to large effects in CP violation and flavor-changing neutral current (FCNC) amplitudes [38]. Particularly, one can generate the enhancement of $C_{\gamma,\ g}^\pm$ at one loop, via intermediate squarks and gluinos, which is due both to the strong-coupling constant and to the removal of chirality suppression present in the SM. Full expressions for the Wilson coefficients generated by gluino exchange at the SUSY scale can be found in Ref. [19]. We are interested here only in the contributions proportional to $1/m_{\tilde{g}}$, which are given by

$$C_{\gamma,\text{SUSY}}^\pm(m_{\tilde{g}}) = \frac{\pi \alpha_s(m_{\tilde{g}})}{m_{\tilde{g}}} [(\delta_{\text{LR}}^{\text{D}})_{21} \pm (\delta_{\text{LR}}^{\text{D}})_{12}^*] F_{\text{SUSY}}(x_{gq}), \tag{34}$$

$$C_{g,\text{SUSY}}^\pm(m_{\tilde{g}}) = \frac{\pi \alpha_s(m_{\tilde{g}})}{m_{\tilde{g}}} [(\delta_{\text{LR}}^{\text{D}})_{21} \pm (\delta_{\text{LR}}^{\text{D}})_{12}^*] G_{\text{SUSY}}(x_{gq}), \tag{35}$$

where $(\delta_{\text{LR}}^{\text{D}})_{ij} = (M_{\text{D}}^2)_{i_{\text{L}} j_{\text{R}}} / m_{\tilde{q}}^2$ denotes the off-diagonal entries of the (down-type) squark mass matrix in the super-CKM basis, $x_{gq} = m_{\tilde{g}}^2 / m_{\tilde{q}}^2$ with $m_{\tilde{g}}$ being the average gluino mass and $m_{\tilde{q}}$ the average squark mass. The explicit expressions of $F_{\text{SUSY}}(x)$ and $G_{\text{SUSY}}(x)$ are given in Ref. [5], but noting that they do not depend strongly on x, it is sufficient, for our purposes, to approximate $F_{\text{SUSY}}(x) \sim F_{\text{SUSY}}(1) = 2/9$ and $G_{\text{SUSY}}(x) \sim G_{\text{SUSY}}(1) = -5/18$. In any case it will be easy to extend the numerology once x_{gq} is better known. Also the determination of the Wilson coefficients in Eqs. (34) and (35) can be improved by the renormalization-group analysis [5,36]. Then by taking $m_{\tilde{g}} = 500$ GeV, $m_t = 174$ GeV, $m_b = 5$ GeV, and $\mu = m_c = 1.25$ GeV, we will have

$$|\text{Im}\, C_\gamma^+|^{\text{SUSY}} = 2.4 \times 10^{-4} \ \ \text{GeV}^{-1} |\text{Im}[(\delta_{\text{LR}}^{\text{D}})_{21} + (\delta_{\text{LR}}^{\text{D}})_{12}^*]|. \tag{36}$$

Using the experimental upper bound on the branching ratio of $K_L \to \pi^0 e^+ e^-$ measured by the KTeV Collaboration [39], the lattice calculation [36] has given

$$|\text{Im}[(\delta_{\text{LR}}^{\text{D}})_{21} + (\delta_{\text{LR}}^{\text{D}})_{12}^*]| < 1.0 \times 10^{-3} \ \ (95\% \ \ \text{C.L.}). \tag{37}$$

Thus from Eqs. (32), (36), and (37), the charge asymmetry in $K^\pm \to \pi^\pm \gamma\gamma$ induced by the EMO in the supersymmetric extensions of the SM could be bound as

$$\left(\frac{\delta\Gamma}{2\Gamma}\right)_{\text{SUSY}}^{\text{EMO}} < \text{a few } \times 10^{-3}, \tag{38}$$

which is significantly larger than the charge asymmetries given in the SM [3,24,28].

In conclusion, we have studied the CP-violating charge asymmetry induced by the electromagnetic penguin operators in $K^\pm \to \pi^\pm \gamma\gamma$ transitions, and supersymmetric extensions of the SM may enhance the Wilson coefficients of these operators, which leads to interesting phenomenology in this study. It is found that the constrain imposed by experiments [39] on the SUSY parameter $|\mathrm{Im}[(\delta^{D}_{LR})_{21}+(\delta^{D}_{LR})^{*}_{12}]|$ allows a significant upper bound on the charge asymmetry given in Eq. (38). Our analysis shows that the charge asymmetry in $K^{\pm}\to\pi^{\pm}\gamma\gamma$ up to 10^{-3} would be a signal of new physics, and thus high-precision measurements of this CP-violating observable might probe interesting extensions of the SM.

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