New physics contribution to $B \to K\pi$ decays in soft collinear effective theory

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We analyze the 5σ difference between the CP asymmetries of the $B^0 \to K^+\pi^-$ and $B^+ \to K^+\pi^0$ decays within the soft collinear effective theory. We find that in the standard model, such a big difference cannot be achieved. We classify then the requirements for the possible New Physics models, which can be responsible for the experimental results. As an example of a New Physics model we study minimal supersymmetric models, and find that the measured asymmetry can be obtained with nonminimal flavor violation.

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I. INTRODUCTION

The first observation of CP violation was in the neutral kaon system in 1964, which was consistent with Cabibbo-Kobayashi-Maskawa (CKM) mechanism and with its simplicity. In the last years, experiments at B-factories have established CP violation in B_d^0 decay. Although the standard model (SM) is able, until now, to account for the CP violating experimental results, CP violation is one of the most interesting aspects and unsolved mysteries of the SM. There are strong hints of additional sources of CP violation beyond the phase in the CKM mixing matrix. The strongest motivation for this suggestion is that the strength of CP violation in the SM is not sufficient to explain the cosmological baryon asymmetry of our Universe. Therefore, it is expected that a sizeable contribution from New Physics (NP) to CP violation in B-meson decays may be probed.

Indeed, there are some discrepancies between the SM expectations and the experimental measurements of the following parameters: $\sin 2\beta_s$ extracted from the mixing CP asymmetry in $B \to J/\psi\phi$ decay [1], $\sin 2\beta$ extracted from the mixing CP asymmetry in $B \to K\phi$ and $B \to K\eta'$ decays [2], and the direct CP asymmetries of $B \to K\pi$ decays. Of these, the $B \to K\pi$ anomaly remains a potential hint for NP that emerges from rare B decays. The current world averages for the branching ratios (BRs) and CP asymmetries of $B \to K\pi$ [2] are summarized in Table I. These results confirm the existence of a nonvanishing difference between the asymmetries of $B^+ \to K^+\pi^0$ and $B^0 \to K^+\pi^-$ beyond 5σ:

$$\mathcal{A}_{CP}(B^+ \to K^+\pi^0) - \mathcal{A}_{CP}(B^0 \to K^+\pi^-) = (14.8 \pm 2.7)\%. \quad (1)$$

It is well known that within the SM, all CP violating processes should be accommodated by the single phase of the CKM, which is the only source of CP violation in the quark sector. This implies tight relations among the CP asymmetries of different processes, which allow stringent tests of the SM, and may therefore lead to the discovery of NP. Indeed, the SM results for the CP asymmetries of $B \to K\pi$, with naive factorization or "improved" Beneke-Buchalla-Neubert-Sachrajda QCD factorization [3] (QCDF), indicate that the above-mentioned two asymmetries are essentially equal [4]. This inconsistency is known as $B \to K\pi$ puzzles and has been considered as a possible hint for physics beyond the SM, with a new source of CP violation. There has been tremendous work over the last few years in order to understand this puzzle of CP asymmetries in $B \to K\pi$ decays.

In this paper, we perform a detailed analysis for the CP asymmetries and branching ratios of $B \to K\pi$ decays in the framework of soft collinear effective theory (SCET) [5,6]. In Ref. [6], the SM contributions to the branching ratios and the CP asymmetries of $B \to K\pi$ have been studied in the frame of the SCET. It was concluded that a small CP asymmetry for $B^+ \to K^+\pi^0$ is predicted and the large discrepancy between the CP asymmetries of $B^0 \to K^+\pi^-$ and $B^+ \to K^+\pi^0$ is difficult to explain in the SM with SCET and a possible new source of New Physics in order to account for these results. Motivated by this conclusion and also by the fact that the difference between these two asymmetries has now reached 5σ, we study the New Physics, in particular, supersymmetry, contributions to these processes and analyze the conditions that may allow for producing the recent experimental results.

The SCET provides a systematic and elegant method for calculating B decays with several relevant energy scales [5–11]. It is based on the fact that the decay of heavy

TABLE I. The latest average results for the BRs and CP asymmetries of $B \to K\pi$ decays.

Decay channel	BR $\times 10^{-6}$	A_{CP}
$K^+\pi^-$	19.4 ± 0.6	-0.098 ± 0.012
$K^+\pi^0$	12.9 ± 0.6	0.050 ± 0.025
$K^0\pi^+$	23.1 ± 1.0	0.009 ± 0.025
$K^0\pi^0$	9.8 ± 0.6	-0.01 ± 0.1

hadrons to highly energetic light hadrons includes three distinct energy scales: the hard energy scale $\sim m_b$, the hard collinear scale $\sim\sqrt{m_b \Lambda_{\rm QCD}}$ and the hadronic soft scale $\sim\Lambda_{\rm QCD}$. Thus, the matching of the weak effective Hamiltonian into the corresponding SCET gauge invariant operators requires two-step matching [11]. First the effective weak Hamiltonian is matched to the corresponding weak Hamiltonian in what is called SCET_I, by integrating out at the hard modes with momentum of order m_b. Second, the SCET_I weak Hamiltonian is matched onto the weak Hamiltonian SCET_{II} by integrating out the hard-collinear modes with $p^2 \sim m_b \Lambda_{\rm QCD}$. Accordingly, the SCET is improving the factorization, obtained from expansion in powers of $\Lambda_{\rm QCD}/m_b$, by generalizing it to allow each of the above-mentioned scales to be considered independently. We will show explicitly that, as in the QCDF approach, the SM results for CP asymmetries of $B \to K\pi$ in SCET are typically not consistent with the observed measurements. This confirms the conclusion that NP is required in order to accommodate the experimental measurements of $B \to K\pi$ CP asymmetries. We will analyze the type of NP needed to resolve the $B \to K\pi$ puzzle and show that it must induce new source of CP violation. As an interesting example of NP, we consider the supersymmetric (SUSY) extension of the SM, using the mass insertion approximation (MIA) in order to perform a model independent analysis.

It is important to note that in order to have significant CP violating effects from SUSY contributions without exceeding the experimental limits of the electric dipole moment (EDM) of electron and neutron, one should consider a SUSY model with nonminimal flavor. In this class of models, like for instance scenarios of nonuniversal trilinear couplings, there are new sources of CP and flavor violation that may lead to significant impacts on the CP asymmetries of $B \to K\pi$, without violating the experimental limits of the electric dipole moment (EDM) of electron or neutron [12]. It has been emphasized in Refs. [4,13] that these phases are crucial in providing a natural explanation for the $B \to K\pi$ puzzle. Indeed, this new source of SUSY CP violating phases induces CP violating phases associated with the electroweak penguins, which are essential with large strong phase in order to resolve the apparent discrepancies between the CP asymmetry of $B^+ \to K^+\pi^0$ and $B^0 \to K^+\pi^-$.

The paper is organized as follows. In Sec. II we discuss the $B \to K\pi$ process in the SCET and present generic expressions for the amplitudes in terms of the Wilson coefficients. Section III is devoted to analyzing the SM contribution to the branching ratios and CP asymmetries of $B \to K\pi$ decays. We show that the branching ratios can be consistent with the experimental data if a large charm penguin contribution is assumed. Nevertheless, the CP asymmetries measurements cannot be accommodated. In Sec. IV we explore the NP effects and possible types of NP that may resolve the puzzle of $B \to K\pi$. We emphasize that a generic feature of any of this NP is that it must introduce a new source of CP violation. In Sec. V we focus our discussion on SUSY extension of the SM. We show that the gluino contribution to the electroweak penguin plays a crucial role in resolving the $B \to K\pi$ puzzle. Finally we summarize our conclusions in Sec. VI.

II. $B \to K\pi$ IN SCET

The full effective weak Hamiltonian $H_{\rm eff}^{\Delta B=1}$ for $\Delta S = 1$ transitions can be expressed via the operator product expansion as

$$H_{\rm eff}^{\Delta B=1} = \frac{G_F}{\sqrt{2}} \sum_{p=u,c} \lambda_p^{(s)} \Big(C_1 Q_1^p + C_2 Q_2^p + \sum_{i=3}^{10} C_i Q_i + C_{7\gamma} Q_{7\gamma} + C_{8g} Q_{8g} \Big) + (Q_i \to \tilde{Q}_i, C_i \to \tilde{C}_i), \quad (2)$$

where $\lambda_p^{(s)} = V_{pb} V_{ps}^{\star}$, with V_{ij} the unitary CKM matrix elements. $C_i \equiv C_i(\mu_b)$ are the Wilson coefficients at low-energy scale $\mu_b \simeq \mathcal{O}(m_b)$. The operators Q_i can be found in Ref. [14]. The operators $Q_{1,2}^p$ refer to the current-current operators, $Q_{3\text{–}6}$ to the QCD penguin operators, and $Q_{7\text{–}10}$ to the electroweak penguin operators, while $Q_{7\gamma}$ and Q_{8g} are the electromagnetic and the chromomagnetic dipole operators, respectively. The operators $\tilde{Q}_i$ are obtained from Q_i by the chirality exchange. It is important to note that the electroweak penguins and the electromagnetic penguin are the only source of isospin violation, which is indicated by the $K\pi$ puzzle.

The calculation of $B \to K\pi$ decays involves the evaluation of the hadronic matrix elements of related operators in the effective Hamiltonian, which is the most uncertain part of this calculation. In the limit in which $m_b \gg \Lambda_{\rm QCD}$ and neglecting QCD corrections in α_s, i.e. in the naive factorization (NF) approach, the hadronic matrix elements of B decays into K and π can be factorized as

$$\langle K\pi|Q_i|B\rangle_{\rm NF} = \langle K|j_1|B\rangle \times \langle \pi|j_2|0\rangle + \langle \pi|j_1|B\rangle \times \langle K|j_2|0\rangle, \quad (3)$$

where $j_{1,2}$ represent bilinear quark currents of local operator Q_i. Therefore, the hadronic matrix element can be usually parametrized by the product of the decay constants and the transition form factors.

In QCDF the hadronic matrix element for $B \to K\pi$ in the heavy quark limit $m_b \gg \Lambda_{\rm QCD}$ can be written as

$$\langle K\pi|Q_i|B\rangle_{\rm QCDF} = \langle K\pi|Q_i|B\rangle_{\rm NF} \times \Big[1 + \sum_n r_n \alpha_s^n + \mathcal{O}\Big(\frac{\Lambda_{\rm QCD}}{m_b}\Big)\Big]. \quad (4)$$

It is clear that in QCDF, the higher order corrections in α_s

break the simple factorization. These corrections can be calculated systematically in terms of short-distance coefficients and meson light-cone distribution functions. However, it turns out that the calculation of the hard spectator interactions and the annihilation amplitude suffer from end-point divergences in this factorization approach. The divergences are parametrized by complex parameters with magnitudes less than one and unconstrained phases. Such parameters are the main source of large theoretical uncertainties in the QCDF mechanism.

The SCET is an interesting framework to study the factorization at hard $\mathcal{O}(m_b)$ and hard-collinear $\mathcal{O}(\sqrt{m_b \Lambda_{\rm QCD}})$ scales. The SCET Lagrangian is obtained at tree level by expanding the full theory Lagrangian in powers of $\lambda = \Lambda_{\rm QCD}/m_b$. This would allow to prove or disprove the factorization to all orders in the strong coupling constant for some B decays into light and energetic particles. Many theoretical works have been done in the SCET, in particular, the matching of QCD $\to$ SCET$_I$ $\to$ SCET$_{II}$ and the derivation of the amplitudes for the B decay into light mesons [5–10]. For $B \to K\pi$, the SCET amplitude can be written as

$$\mathcal{A}^{\rm SCET}_{B\to K\pi} = -i\langle K\pi | H^{\rm SCET}_{\rm eff} | B\rangle = \mathcal{A}^{\rm LO}_{B\to K\pi} + \mathcal{A}^{\chi}_{B\to K\pi} + \mathcal{A}^{\rm ann}_{B\to K\pi} + \mathcal{A}^{c.c}_{B\to K\pi} \tag{5}$$

where $\mathcal{A}^{\rm LO}_{B\to K\pi}$ denotes the leading-order amplitude in the expansion $1/m_b$ (including correction of order α_s), $\mathcal{A}^{\chi}_{B\to K\pi}$ denotes the chirally enhanced penguin amplitude, $\mathcal{A}^{\rm ann}_{B\to K\pi}$ denotes the annihilation amplitude and $\mathcal{A}^{c.c}_{B\to K\pi}$ denotes the long- distance charm penguin contributions.

The leading-order amplitude, $\mathcal{A}^{\rm LO}_{B\to K\pi}$, is given by

$$\mathcal{A}^{\rm LO}_{B\to K\pi} = \frac{G_F m_B^2}{\sqrt{2}}\bigg[f_K\bigg(\int_0^1 du dz T_{KJ}(u,z)\zeta_J^{B\pi}(z)\phi_K(u) + \zeta^{B\pi}\int_0^1 du T_{K\zeta}(u)\phi_K(u)\bigg) + (K\leftrightarrow\pi)\bigg]. \tag{6}$$

The hard kernels $T_{(K,\pi)\zeta}$ and $T_{(K,\pi)J}$ are calculable in terms of the Wilson coefficients C_i and can be found in Ref. [15]. The parameters $\zeta^{B(K,\pi)}$, $\zeta_J^{B(K,\pi)}$ are treated as hadronic parameters that can be determined through the fit to the nonleptonic decay data. The current data can be used to determine $\zeta^{B\pi}$, $\zeta_J^{B\pi}$. However these data are not sufficient to determine ζ^{BK} and ζ_J^{BK} and hence we assume $\zeta_J^{BK} = \zeta_J^{B\pi}$ and $\zeta^{BK} = \zeta^{B\pi}$ in the limit of exact $SU(3)$. One may expect about 10%–20% deviation in the values of these parameters in case of $SU(3)$ breaking.

It is important to note that as long as the logarithms of the ratios of the hard scale (m_b) to the soft-collinear (Λm_b) and soft (Λ) scales are not resummed, the QCDF and SCET factorization formulas are identical. Therefore, Eq. (6) for the expression of $A^{\rm LO}_{B\to K\pi}$ includes as well the end-point singular contribution mentioned above in QCDF scheme. In fact, the form factors $\xi_J^{B\pi}(z)$, which are extracted from the data, could be expressed as end-point singular convolutions between the pion and B-meson light-cone wave functions.

Chiraly enhanced penguins amplitude $\mathcal{A}^{\chi}_{B\to K\pi}$ is generated through including corrections of order $\alpha_s(\mu_h)\times(\mu_M\Lambda/m_b^2)$ where μ_M is the chiral scale parameter. μ_M is defined as the ratio of the squared meson mass to the sum of its constituent quark masses. For kaons and pions $\mu_M \sim \mathcal{O}(2)$ GeV and hence chiraly enhanced terms can compete with the order $\alpha_s(\mu_h)(\Lambda/m_b)$ terms. The chiraly enhanced amplitude for $B\to K\pi$ decays is given by

$$A^{\chi}_{\bar{B}\to K\pi} = \frac{G_F m_B^2}{\sqrt{2}}\bigg\{ -\frac{\mu_K f_K}{3m_B}\zeta^{B\pi}\int_0^1 du R_K(u)\phi^K_{pp}(u) - \frac{\mu_K f_K}{3m_B}\int_0^1 dudz R^J_K(u,z)\zeta_J^{B\pi}(z)\phi^K_{pp}(u) - \frac{\mu_\pi f_K}{6m_B}\int_0^1 dudz R^\chi_K(u,z)\zeta_\chi^{B\pi}(z)\phi^K(u) + (K\leftrightarrow\pi)\bigg\}. \tag{7}$$

The hard kernels R_K, R_π, R^J_K, R^J_π, R^χ_K and R^χ_π depend also on C_i, as shown in [10].

Annihilation amplitudes $\mathcal{A}^{\rm ann}_{B\to K\pi}$ have been studied in Refs. [16–19]. In the framework of SCET, the annihilation contribution becomes factorizable and real at leading order, $\mathcal{O}(\alpha_s(m_b)\Lambda/m_b)$. Complex annihilation contributions may occur at higher order, $\mathcal{O}(\alpha_s^2(\sqrt{m_b\Lambda})\Lambda/m_b)$ [20]. In our numerical analysis, we will not include the contributions from penguin annihilations, since they are real, at the order we consider, and are quite small with large uncertainty [10,20]. It is worth mentioning that there are some question marks related to the SCET result for $A^{\rm ann}_{B\to K\pi}$. It is expected that the approach adopted in computing the LO expression may lead to a divergent annihilation contribution, which therefore requires a reintroduction of complex parameter as in QCDF. This discussion is beyond the scope of this paper, especially in the case of neglecting the annihilation amplitude.

The long-distance charm penguin amplitude $\mathcal{A}^{c.c}_{B\to K\pi}$ is given as follows:

$$\mathcal{A}^{c.c}_{B\to K\pi} = |\mathcal{A}^{c.c}_{B\to K\pi}| e^{i\delta_{cc}} \tag{8}$$

where δ_{cc} is the strong phase of the charm penguin. The modulus and the phase of the charm are fixed, through the fitting with nonleptonic decays, namely $B\to\pi\pi$, assuming $\mathcal{A}^{c.c}_{B\to K\pi} = \mathcal{A}^{c.c}_{B\to\pi\pi}$, as follows [21]:

$$|A_{c.c}| = (46 \pm 0.8)\times 10^{-4}, \qquad \delta_{c.c} = 156^o \pm 6^o. \tag{9}$$

The charm penguin can be considered as one of the main differences between SCET and QCDF. In QCDF, it is factorized in the limit of $1/m_b$. However, in SCET, since $m_c \sim m_b/2$ there may be configurations, where the charm penguin implies a long-distance effect. Thus, it has been parametrized and fitted from the data. It is also worth noting that in SCET the charm penguin is the main source of strong phases in the decay amplitudes. All strong phases for other terms vanish at the leading order.

The unitarity of the CKM matrix allows to write the amplitude of any B-decay as $A = \lambda_u^{(f)} A_u + \lambda_c^{(f)} A_c$, where $\lambda_p^{(f)} = V_{pb}^* V_{pf}$. Thus, one can generally parametrize the contributions to the amplitudes of $B \to K\pi$ as follows:

$$\begin{aligned} A(B^+ \to K^0\pi^+) &= \lambda_u A + \lambda_c P, \\ \sqrt{2}A(B^+ \to K^+\pi^0) &= \lambda_u(T + C + A) + \lambda_c(P + P_{\rm EW}), \\ A(B^0 \to K^+\pi^-) &= \lambda_u T + \lambda_c(P + P_{\rm EW}^C), \\ \sqrt{2}A(B^0 \to K^0\pi^0) &= \lambda_u C - \lambda_c(P - P_{\rm EW} + P_{\rm EW}^C), \end{aligned} \tag{10}$$

where the real parameters, T, C, A, P, $P_{\rm EW}$, and $P_{\rm EW}^C$, represent a colored allowed tree, a color-suppressed tree, annihilation, QCD penguin, electroweak penguin, and suppressed electroweak penguin diagrams, respectively. The four $B \to K\pi$ decay amplitudes are related by the following isospin relation:

$$\sqrt{2}A(B^0 \to K^0\pi^0) + A(B^+ \to K^0\pi^+) - \sqrt{2}A(B^+ \to K^+\pi^0) + A(B^0 \to K^+\pi^-) = 0. \tag{11}$$

The explicit dependence of these parameters on the corresponding Wilson coefficients can be found in Refs. [5,7–10]. Fixing the experimental inputs and the SM parameters to their center values, one finds the following dependence of these parameters on the Wilson coefficients (at NLO in α_s expansion of SCET + $SU(3)$ flavor symmetry):

$$\begin{aligned} \hat{A} &= -(0.0003 + 0.0005i)C_1 - 0.0134C_{10} + (0.0233 - 0.0009i)C_3 + 0.0268C_4 + 0.0113C_5 + 0.034C_6 - 0.0057C_7 \\ &\quad - 0.017C_8 - (0.012 - 0.0005i)C_9 - 0.0009C_{8g}, \\ \hat{P} &= (-0.0004 - 0.0003i)C_1 - 0.013C_{10} + (0.0234 - 0.0009i)C_3 + 0.027C_4 + 0.0113C_5 + 0.034C_6 - 0.006C_7 \\ &\quad - 0.017C_8 - (0.012 - 0.0005i)C_9 - 0.0009C_{8g} - (0.004 - 0.002i), \\ \hat{P}_{\rm EW}^C &= 0.017C_7 + 0.051C_8 + (0.035 - 0.0014i)C_9 + 0.04C_{10}, \\ \hat{P}_{\rm EW} &= -0.016C_7 + (0.056 - 0.0014i)C_8 + (0.068 - 0.0014i)C_9 + (0.066 - 0.0013i)C_{10}, \\ \hat{C} &= (0.017 - 0.0004i)C_1 + (0.039 - 0.001i)C_{10} + 0.022C_2 - (0.023 - 0.0009i)C_3 - 0.027C_4 - 0.011C_5 \\ &\quad - 0.034C_6 - 0.027C_7 + (0.022 - 0.0014i)C_8 + 0.0009C_{8g} + (0.045 - 0.0005i)C_9, \\ \hat{T} &= (0.027 - 0.001i)C_1 + 0.027C_{10} + (0.023 - 0.001i)C_2 + (0.024 - 0.001i)C_3 + 0.027C_4 + 0.011C_5 + 0.034C_6 \\ &\quad + 0.0113C_7 + 0.034C_8 - 0.0009C_{8g} + (0.024 - 0.001i)C_9, \end{aligned} \tag{12}$$

where $\hat{X}$ is defined as $\hat{X} = \sqrt{2}X/G_F m_B^2$ with $X \equiv A, T, C, P, P_{\rm EW}, P_{\rm EW}^C$. The above results correspond to the total amplitudes including the chirally enhanced penguin with inclusion of the charm penguin as a nonperturbative contribution fitted from the experimental data. Note that the charm penguin contributes only to the QCD penguin P, and it is fixed from the data of $B \to \pi\pi$ processes.

Here a few comments are in order: (i) At leading order, the only source of the strong phases is the charm penguin, however at next-to-leading-order correction, small strong phases may emerge. (ii) In the combined SCET + $SU(3)$, one finds that $C \sim T$, hence there is no color suppression. (iii) There is no undetermined strong phase in the amplitudes T, C, $P_{\rm EW}$, $P_{\rm EW}^C$, unlike the QCDF. Thus, the relative sign of CP asymmetries is predicted. (iv) The amplitudes $P_{\rm EW}$ and $P_{\rm EW}^C$ receive contributions through the electroweak penguin operators O_{7-10}. Unlike the gluonic penguins, the electroweak (γ- and Z- mediated) penguins distinguish the up from the down quark pairs in the final state. Therefore, if they are not suppressed, they may account for the difference between the CP asymmetries in the two isospin related decays of Eq. (1).

III. SM CONTRIBUTION TO THE CP ASYMMETRY OF $B \to K\pi$

In this section we reappraise the SM predictions for the CP asymmetries of $B \to K\pi$ decays in SCET [6]. In the NDR scheme taking $\alpha_s(m_Z) = 0.118$, $m_t = 170.9$ GeV, and $m_b = 4.7$ GeV, the Wilson coefficients are given by

$$\begin{aligned} C_{1-10}(m_b) &= \{1.078, -.177, .012, -.0335, .0095, -.040, \\ &\qquad 1 \times 10^{-4}, 4.2 \times 10^{-4}, -9.7 \times 10^{-3}, \\ &\qquad 1.9 \times 10^{-3}\}, \\ C_{7\gamma}(m_b) &= -.316, \qquad C_{8g}(m_b) = -0.149. \end{aligned} \tag{13}$$

As can be seen from these values, the SM contributions to the electroweak penguins C_7–C_{10} are quite suppressed. Thus, one expects that the EW penguins in the SM are negligible, hence the $B \to K\pi$ asymmetries are dominated by the QCD penguins, which give universal contributions to the four-decay channel. Accordingly, it is expected that the SM results for the CP asymmetries of different $B \to K\pi$ channels are very close. Since the SM Wilson coefficients are real, one can rewrite the amplitude of $B \to K\pi$

in Eq. (10) as

$$\begin{aligned} A(B^+ \to K^0\pi^+) &= \lambda_c^{(s)} P[1 + r_A e^{i(\delta_A - \gamma)}], \\ A(B^0 \to K^+\pi^-) &= \lambda_c^{(s)} P[1 + (r_{\rm EW}^C e^{i\delta_{\rm EW}^C} + r_T e^{i(\delta_T - \gamma)})], \\ \sqrt{2}A(B^+ \to K^+\pi^0) &= \lambda_c^{(s)} P[1 + (r_{\rm EW} e^{i\delta^{\rm EW}} + r_T e^{i(\delta_T - \gamma)} \\ &\quad + r_C e^{i(\delta_C - \gamma)} + r_A e^{i(\delta_A - \gamma)})], \\ \sqrt{2}A(B^0 \to K^0\pi^0) &= \lambda_c^{(s)} P[-1 + (r_{\rm EW} e^{i\delta_{\rm EW}} - r_{\rm EW}^C e^{i\delta_{\rm EW}^C} \\ &\quad + r_C e^{i(\delta_C - \gamma)})], \end{aligned} \tag{14}$$

where the parameters δ_J, with J stands for T, C, A, EW, EW^C, are the CP conserving (strong) phase and r_J are defined as

$$\begin{aligned} r_T e^{i\delta_T} &= \left|\frac{\lambda_u^{(s)}}{\lambda_c^{(s)}}\right| \frac{T}{P}, \qquad r_C e^{i\delta_C} = \left|\frac{\lambda_u^{(s)}}{\lambda_c^{(s)}}\right| \frac{C}{P}, \\ r_A e^{i\delta_A} &= \left|\frac{\lambda_u^{(s)}}{\lambda_c^{(s)}}\right| \frac{A}{P}, \qquad r_{\rm EW} e^{i\delta_{\rm EW}} = \frac{P_{\rm EW}}{P}, \\ r_{\rm EW}^C e^{i\delta_{\rm EW}^C} &= \frac{P_{\rm EW}^C}{P}. \end{aligned} \tag{15}$$

As can be seen from Eq. (10), P is dominated by the large charm penguin. Therefore, one finds that all the above ratios are quite suppressed and also have one single strong phase, which is essentially $\delta_{c.c}$. Namely, one obtains the following results:

$$\begin{aligned} r_T e^{i\delta_T} &= 0.06 e^{-2.91i}, \qquad r_C e^{i\delta_C} = 0.05 e^{-2.92i}, \\ r_A e^{i\delta_A} &= 0.006 e^{0.54i}, \qquad r_{\rm EW} e^{i\delta_{\rm EW}} = 0.08 e^{0.23i}, \\ r_{\rm EW}^C e^{i\delta_{\rm EW}^C} &= 0.04 e^{0.2i}. \end{aligned} \tag{16}$$

From these results, one notices that in SCET the ratio between the color-suppressed tree and color-allowed tree is enhanced, so $|C/T| \sim 1$, unlike the corresponding ratio in QCDF. This enhancement is due to the suppression of T, not because enhancement of C. In this approach, one finds $r_T \sim r_C$ and $r_{\rm EW} \sim r_{\rm EW}^C$, which means there is no color suppression. However, even if the color-suppressed tree and electroweak penguin $(C, P_{\rm EW}^C)$ are enhanced and become of the order of the color-allowed tree and electroweak penguin $(T, P_{\rm EW})$, it is not possible to resolve the puzzle $B \to K\pi$ CP asymmetry in the framework of the SM, due to a lack of CP violation as emphasized in Ref. [4]. Because of the dominance of $A_{c.c.}$ in P, hence $r_J \ll 1$, the following relation between the amplitudes of different channels is established:

$$A_{K^0\pi^+} \simeq A_{K^+\pi^-} \simeq \sqrt{2}A_{K^+\pi^0} \simeq \sqrt{2}A_{K^0\pi^0}. \tag{17}$$

The branching ratio of $B \to K\pi$ is given by

$$\mathrm{BR}\,(B \to K\pi) = \frac{1}{\Gamma_{\rm tot}} \frac{[(M_B^2 - (m_K + m_\pi)^2)(M_B^2 - (m_K - m_\pi)^2)]^{1/2}}{16\pi M_B^3}[|A_{K\pi}|^2 + |\bar{A}_{K\pi}|^2]. \tag{18}$$

Therefore, the BRs also satisfy the relation:

$$\mathrm{BR}_{K^0\pi^+} \simeq \mathrm{BR}_{K^+\pi^-} \simeq 2\mathrm{BR}_{K^+\pi^0} \simeq 2\mathrm{BR}_{K^0\pi^0}, \tag{19}$$

which is consistent with the data given in Table I. However, the magnitude of the BR is sensitive to the value of P and hence to the value of the charm penguin $A_{c.c}$. In fact, for negligible charm penguin, i.e., $A_{c.c} = 0$ one finds that BRs are given by

$$\begin{aligned} \mathrm{BR}_{K^0\pi^+} &= 2.1 \times 10^{-6}, \qquad \mathrm{BR}_{K^+\pi^-} = 2.3 \times 10^{-6}, \\ \mathrm{BR}_{K^+\pi^0} &= 1.4 \times 10^{-6}, \qquad \mathrm{BR}_{K^0\pi^0} = 0.9 \times 10^{-6}. \end{aligned} \tag{20}$$

These results are smaller than the experimental measurements. Therefore, it is appealing that the large charm penguin is essential for the consistency of the SCET. For the value of $A_{c.c.}$ in Eq. (9), one finds significant enhancement for the BRs and they become close to the experimental results, namely, they are now given by

$$\begin{aligned} \mathrm{BR}_{K^0\pi^+} &= 20.5 \times 10^{-6}, \qquad \mathrm{BR}_{K^+\pi^-} = 21.1 \times 10^{-6}, \\ \mathrm{BR}_{K^+\pi^0} &= 11.2 \times 10^{-6}, \qquad \mathrm{BR}_{K^0\pi^0} = 9.7 \times 10^{-6}. \end{aligned} \tag{21}$$

In order to understand the dependence of the CP asymmetries on different contributions, we will neglect small r_J^2 corrections. However our numerical results are based on the complete expressions of the asymmetries, which turn out to be quite close to the approximated ones. Keeping linear terms in r_J, one finds that the $B \to K\pi$ CP asymmetries can be written as

$$\begin{aligned} A^{CP}_{B^+\to K^0\pi^+} &= \frac{2r_A \sin\delta_A \sin\gamma}{1 + 2r_A \cos\delta_A \cos\gamma}, \\ A^{CP}_{\bar{B}^0\to K^+\pi^-} &= \frac{2r_T \sin\delta_T \sin\gamma}{1 + 2r^C_{\rm EW} \cos\delta^C_{\rm EW} + 2r_T \cos\delta_T \cos\gamma}, \\ A^{CP}_{B^+\to K^+\pi^0} &= \frac{2r_T \sin\delta_T \sin\gamma + 2r_C \sin\delta_C \sin\gamma + 2r_A \sin\delta_A \sin\gamma}{1 + 2r_{\rm EW} \cos\delta_{\rm EW} + 2r_C \cos\delta_C \cos\gamma + 2r_T \cos\delta_T \cos\gamma + 2r_A \cos\delta_A \cos\gamma}, \\ A^{CP}_{B^0\to K^0\pi^0} &= \frac{-2r_C \sin\delta_C \sin\gamma}{1 - 2r_{\rm EW} \cos\delta_{\rm EW} + 2r^C_{\rm EW} \cos\delta^C_{\rm EW} - 2r_C \cos\delta_C \cos\gamma}. \end{aligned} \tag{22}$$

It is interesting to note that without charm penguin contribution, although r_J is not suppressed, all the *CP* asymmetries of $B \to K\pi$ decays are quite small, $\mathcal{O}(0.01)$, which is not consistent with the experimental results reported above in Table I. This is due to the lack of large strong phases. As mentioned, the charm penguin in SCET is the main source of strong phases. Therefore these phases associated with r_J are essentially given by $\pm 1/P$. This can be checked in Eq. (15), where one observes the following relation:

$$\sin\delta_T = \sin\delta_C = \sin\delta_A = -\sin\delta_{\rm EW} = -\sin\delta^C_{\rm EW} = -\sin\delta_P. \tag{23}$$

It is now clear that the above expression of the *CP* asymmetries cannot lead to $A^{CP}_{K^+\pi^-}$ and $A^{CP}_{K^+\pi^0}$ with different signs. In fact, one can approximate these two asymmetries as follows: $A^{CP}_{K^+\pi^-} \simeq 2r_T \sin\delta \sin\gamma$ and $A^{CP}_{K^+\pi^0} \simeq 2(r_T + r_C)\sin\delta \sin\gamma$, which lies between $A^{CP}_{K^+\pi^-}$ and $2A_{K^+\pi^-}$. One can check this conclusion numerically. For instance, with a charm penguin fixed by $B \to \pi\pi$ [6], one finds the following asymmetries:

$$A^{CP}_{B^+\to K^0\pi^+} = -0.01, \qquad A^{CP}_{B^0\to K^+\pi^-} = -0.03, \qquad A^{CP}_{B^+\to K^+\pi^0} = -0.04, \qquad A^{CP}_{B^0\to K^0\pi^0} = 0.02. \tag{24}$$

Note that the EW penguins violate the isospin symmetry, hence they are natural candidates for explaining the discrepancy between $A^{CP}_{K^+\pi^-}$ and $A^{CP}_{K^+\pi^0}$. However, as we have seen, within the SM, these two asymmetries are not sensitive to the values of $r_{\rm EW}$ and $r^c_{\rm EW}$. This is due to the fact that the EW penguins are real in the SM and hence they have no interference with the QCD penguin P. As emphasized in Ref. [13], a possible solution for the $B \to K\pi$ puzzle is to have a new source of *CP* violation that generates *CP* phases for the EW penguins. This possibility can be implemented in supersymmetric models and has been checked within the framework of QCDF in Refs. [4,13].

IV. NEW PHYSICS EFFECTS AND *CP* ASYMMETRIES OF $B \to K\pi$ IN SCET

In this section we analyze the type of general NP beyond the SM that can account for the *CP* asymmetries of $B \to K\pi$ and explain the discrepancy between $A^{CP}_{K^+\pi^-}$ and $A^{CP}_{K^+\pi^0}$. As mentioned above and discussed in detail in Ref. [4], this NP must contain a new source of *CP* violation beyond the CKM phase. The impact of any NP beyond the SM appears only in the Wilson coefficients at electroweak scale. Therefore, the total Wilson coefficients can be written as

$$C_i = C_i^{\rm SM} + C_i^{\rm NP}, \qquad i = 1, \ldots, 10, 7\gamma, 8g, \tag{25}$$

where $C_i^{\rm NP}$ are generally complex, i.e., they have a *CP* violating phase, unlike the $C_i^{\rm SM}$. Also the NP is expected to have relevant contributions to the penguins and not to the tree processes, which are dominated by the SM effects. Therefore, one can assume that the color-tree and color-suppressed parameters remain as in the SM, i.e., $T = T^{\rm SM}$ and $C = C^{\rm SM}$, while the penguin parameters are given by

$$\begin{aligned} Pe^{i\theta_P}e^{i\delta_P} &= |P^{\rm SM}|e^{i\delta_{c.c}} + |P^{\rm NP}|e^{i\phi_P} \\ &= |P^{\rm SM}|[e^{i\delta_{c.c}} + \kappa_P e^{i\phi_P}], \end{aligned} \tag{26}$$

$$\begin{aligned} P_{\rm EW}e^{i\theta_{\rm EW}} &= |P^{\rm SM}_{\rm EW}| + |P^{\rm NP}_{\rm EW}|e^{i\phi_{\rm EW}} \\ &= |P^{\rm SM}_{\rm EW}|[1 + \kappa_{\rm EW}e^{i\phi_{\rm EW}}], \end{aligned} \tag{27}$$

$$\begin{aligned} P_{{\rm EW}^C}e^{i\theta^C_{\rm EW}} &= |P^{\rm SM}_{{\rm EW}^C}| + |P^{\rm NP}_{{\rm EW}^C}|e^{i\phi_{{\rm EW}^C}} \\ &= |P^{\rm SM}_{{\rm EW}^C}|[1 + \kappa^C_{\rm EW}e^{i\phi^C_{\rm EW}}]. \end{aligned} \tag{28}$$

Here we assume that the only source of strong phase is $\delta_{c.c}$ in $P^{\rm SM}$. As mentioned in the previous section, a large charm penguin contribution is very crucial in the SCET in order to get the branching ratio of $B \to K\pi$ decays consistent with the experimental measurements. Furthermore, it is also needed to allow for a large strong phase, which is crucial for generating a large *CP* asymmetry. In order to generalize the parametrization of $B \to K\pi$ in Eq. (14), one should rewrite P as $P = |P|e^{i\delta_P}e^{i\theta_P}$, where δ_P and θ_P are the strong (*CP* conserving) and *CP* violating phases associated with P, which can be determined as follows:

$$\begin{aligned} \delta_P &= \tan^{-1}\left(\frac{\sin\delta_{c.c}}{\cos\delta_{c.c} + \kappa_P \cos\phi_P}\right), \\ \theta_P &= \tan^{-1}\left(\frac{\kappa_P \sin\phi_P}{\kappa_P \cos\phi_P + \cos\delta_{c.c}}\right). \end{aligned} \tag{29}$$

Similarly, $\theta_{\rm EW}$ and $\theta^C_{\rm EW}$ can be defined in terms of $\phi_{\rm EW}$ and $\phi^C_{\rm EW}$. In this case, the ratio between the EW and QCD penguins can be written as

$$\frac{P_{\rm EW}}{P} = r_{\rm EW} e^{-i\delta_P} e^{i(\theta_{\rm EW}-\theta_P)}, \qquad \frac{P^C_{\rm EW}}{P} = r^C_{\rm EW} e^{-i\delta_P} e^{i(\theta^C_{\rm EW}-\theta_P)}, \tag{30}$$

where $r_{\rm EW}$ and $r^C_{\rm EW}$ are given by

$$r_{\rm EW} = (r_{\rm EW})^{\rm SM} \left| \frac{1+\kappa_{\rm EW} e^{i\phi_{\rm EW}}}{1+\kappa_P e^{i(\phi_P-\delta_{c.c})}} \right|, \qquad r^C_{\rm EW} = (r^C_{\rm EW})^{\rm SM} \left| \frac{1+\kappa^C_{\rm EW} e^{i\phi^C_{\rm EW}}}{1+\kappa_P e^{i(\phi_P-\delta_{c.c})}} \right|. \tag{31}$$

Note that the strong phases still satisfy the relation in Eq. (23), as in the SM. This leads to the following parametrization for the $B \to K\pi$ amplitudes:

$$\begin{aligned} A(B^+ \to K^0\pi^+) &= \lambda_c^{(s)} P[e^{i\theta_P} + r_A e^{i(\delta_A-\gamma)}], \\ A(B^0 \to K^+\pi^-) &= \lambda_c^{(s)} P[e^{i\theta_P} + (r^C_{\rm EW} e^{i(\theta^C_{\rm EW}+\delta^C_{\rm EW})} + r_T e^{i(\delta_T-\gamma)})], \\ \sqrt{2}A(B^+ \to K^+\pi^0) &= \lambda_c^{(s)} P[e^{i\theta_P} + (r_{\rm EW} e^{i(\theta_{\rm EW}+\delta_{\rm EW})} + r_T e^{i(\delta_T-\gamma)} + r_C e^{i(\delta_C-\gamma)} + r_A e^{i(\delta_A-\gamma)})], \\ \sqrt{2}A(B^0 \to K^0\pi^0) &= \lambda_c^{(s)} P[-e^{i\theta_P} + (r_{\rm EW} e^{i(\theta_{\rm EW}+\delta_{\rm EW})} - r^C_{\rm EW} e^{i(\theta^C_{\rm EW}+\delta^C_{\rm EW})} + r_C e^{i(\delta_C-\gamma)})]. \end{aligned} \tag{32}$$

In this case, one finds that the approximate expressions for the CP asymmetries in Eq. (22) can be generalized as follows:

$$\begin{aligned} A^{CP}_{B^+\to K^0\pi^+} &= \frac{2r_A \sin\delta_A \sin(\theta_P+\gamma)}{1+2r_A\cos\delta_A\cos(\theta_P+\gamma)}, \\ A^{CP}_{\bar{B}^0\to K^+\pi^-} &= \frac{2r_T\sin\delta_T\sin(\theta_P+\gamma) + 2r^C_{\rm EW}\sin\delta^C_{\rm EW}\sin(\theta_P-\theta^C_{\rm EW})}{1+2r_T\cos\delta_T\cos(\theta_P+\gamma)+2r^C_{\rm EW}\cos\delta^C_{\rm EW}\cos(\theta_P-\theta^C_{\rm EW})}, \\ A^{CP}_{B^+\to K^+\pi^0} &= \frac{2r_{\rm EW}\sin\delta_{\rm EW}\sin(\theta_P-\theta_{\rm EW}) + 2[r_T\sin\delta_T + r_C\sin\delta_C + r_A\sin\delta_A]\sin(\theta_P+\gamma)}{1+2r_{\rm EW}\cos\delta_{\rm EW}\cos(\theta_P-\theta_{\rm EW}) + 2[r_T\cos\delta_T+r_C\cos\delta_C+r_A\cos\delta_A]\cos(\theta_P+\gamma)}, \\ A^{CP}_{B^0\to K^0\pi^0} &= \frac{-2r_{\rm EW}\sin\delta_{\rm EW}\sin(\theta_P-\theta_{\rm EW}) + 2r^C_{\rm EW}\sin\delta^C_{\rm EW}\sin(\theta_P-\theta^C_{\rm EW}) - 2r_C\sin\delta_C\sin(\theta_P+\gamma)}{1-2r_{\rm EW}\cos\delta_{\rm EW}\cos(\theta_P-\theta_{\rm EW}) + 2r^C_{\rm EW}\cos\delta^C_{\rm EW}\cos(\theta_P-\theta^C_{\rm EW}) - 2r_C\cos\delta_C\cos\theta_P+\gamma}. \end{aligned} \tag{33}$$

If one assumes $r_C \sim r_T$, and neglects the small r_A, then the CP asymmetries $A^{CP}_{K^+\pi^-}$ and $A^{CP}_{K^+\pi^0}$, which are not consistent with the SM results, can be written as:

$$\begin{aligned} A^{CP}_{K^+\pi^-} &\simeq \frac{2\sin\delta_P[-r_T\sin(\theta_P+\gamma) + r^C_{\rm EW}\sin(\theta_P-\theta^C_{\rm EW})]}{1+2r_T\cos\delta_P\cos(\theta_P+\gamma)+2r^C_{\rm EW}\cos\delta_P\cos(\theta_P-\theta^C_{\rm EW})}, \\ A^{CP}_{K^+\pi^0} &\simeq \frac{2\sin\delta_P[r_{\rm EW}\sin(\theta_P-\theta_{\rm EW}) - 2r_T\sin(\theta_P+\gamma)]}{1+2r_{\rm EW}\cos\delta_P\cos(\theta_P-\theta_{\rm EW})+4r_T\cos\delta_P\cos(\theta_P+\gamma)}. \end{aligned} \tag{34}$$

Therefore, the difference between these two asymmetries is now given by

$$A^{CP}_{K^+\pi^0} - A^{CP}_{K^+\pi^-} \simeq 2\sin\delta_P[r_{\rm EW}\sin(\theta_P-\theta_{\rm EW}) - r_T\sin(\theta_P+\gamma) - r^C_{\rm EW}\sin(\theta_P-\theta^C_{\rm EW})]. \tag{35}$$

Note that the denominators in Eq. (34) can be approximated to one if large phases are considered to maximize the asymmetries. According to Eq. (1), this difference should be of order $\mathcal{O}(0.14)$ in order to match the current experimental results. Thus one finds

$$r_{\rm EW}\sin(\theta_P-\theta_{\rm EW}) - r_T\sin(\theta_P+\gamma) - r^C_{\rm EW}\sin(\theta_P-\theta^C_{\rm EW}) \simeq \frac{0.07}{\sin\delta_P}. \tag{36}$$

Moreover, the result of $A^{CP}_{K^+\pi^-}$ implies that

$$-r_T\sin(\theta_P+\gamma) + r^C_{\rm EW}\sin(\theta_P-\theta^C_{\rm EW}) \sim \frac{-0.049}{\sin\delta_P}. \tag{37}$$

From these relations, one gets:

$$r_{\rm EW}\sin(\theta_P-\theta_{\rm EW}) - 2r^C_{\rm EW}\sin(\theta_P-\theta^C_{\rm EW}) \simeq \frac{0.12}{\sin\delta_P}. \tag{38}$$

This condition can be fulfilled if one of the following scenarios takes place:

(i) $r_{\rm EW}\sin(\theta_P-\theta_{\rm EW}) \sim 0.12/\sin\delta_P$, while $r^C_{\rm EW}\sin(\theta_P-\theta^C_{\rm EW}) \lesssim \mathcal{O}(0.01)$, which could be due to smallness of $r^C_{\rm EW}$ or $\theta_P \sim \theta^C_{\rm EW}$. Note that if $\delta_P \sim \delta_{c.c}$, then $r_{\rm EW}\sin(\theta_P - \theta_{\rm EW}) \sim 0.3$. In this case, the required NP should enhance the value of $r_{\rm EW}$ to be larger than $|0.12/\sin\delta_P|$ and induce CP violating phases such that $\sin(\theta_P - \theta_{\rm EW}) \sim \mathcal{O}(1)$, i.e., $\theta_{\rm EW} \simeq \theta_P - \pi/2$. The phase θ_P can be fixed from $A^{CP}_{K^+\pi^-}$ which in this scenario is given by $2r_T \sin\delta_P \sin(\theta_P + \gamma)$.

(ii) $r_{\rm EW} \sim r^C_{\rm EW}$ and $\theta_{\rm EW} \sim \theta^C_{\rm EW}$. In this case, the required NP should lead to $r_{\rm EW}\sin(\theta_P - \theta_{\rm EW}) \sim r^C_{\rm EW}\sin(\theta_P - \theta^C_{\rm EW}) \sim -0.12/\sin\delta_P$. Therefore, $r_{\rm EW}$ should also be larger than $|0.12/\sin\delta_P|$ and $\sin(\theta_P - \theta^C_{\rm EW}) \sim \mathcal{O}(-1)$, i.e., $\theta_{\rm EW} \sim \theta^C_{\rm EW} \sim \theta_P + \pi/2$.

(iii) Another possibility is that $r_{\rm EW}\sin(\theta_P - \theta_{\rm EW}) \lesssim \mathcal{O}(0.01)$ and $r^c_{\rm EW}\sin(\theta_P - \theta_{{\rm EW}^c}) \sim -0.06/\sin\delta_P$. It may be natural to think that color-allowed contributions should dominate the color-suppressed ones, therefore this scenario requires a NP that implies: $\theta_P \sim \theta_{\rm EW}$ and $\sin(\theta_P - \theta^C_{\rm EW}) \sim -0.06/(r^C_{\rm EW}\sin\delta_P)$.

It is important to note that in these three marked scenarios, the new CP violating phases are crucial and play an important role in modifying the $B \to K\pi$ CP asymmetries and moving them towards the experimental measurements. This could be an interesting test for the correct NP that we should consider as extension of the SM. In the next section we will check the possibility that SUSY can resolve the puzzle of $B \to K\pi$ as it can do in the QCDF [4], and if it is so, which scenario of the above three can be implemented in SUSY models. It is also worth mentioning that if the denominators of Eq. (33) are less than one, then the value of the CP asymmetries can be enhanced and smaller values of CP phases could be sufficient for accommodating the experimental results of CP asymmetries of $B \to K\pi$ decays.

Before concluding this section, it is worth mentioning that in QCDF there is more than one source of strong phases, therefore one may adjust the sign of $\delta_{\rm EW}$ and $\delta^C_{\rm EW}$ such that the difference between $A^{CP}_{K^+\pi^-}$ and $A^{CP}_{K^+\pi^0}$ can be obtained without any tight relation between the CP violating phases of the QCD and EW penguins, like those obtained in SCET. Accordingly, it is expected to be more difficult for NP to account for the CP asymmetry of $B \to K\pi$ decays in SCET than in other frames of hadron dynamics.

V. SUSY CONTRIBUTIONS TO THE *CP* ASYMMETRY OF $B \to K\pi$ IN SCET

Now, we consider SUSY as a potential candidate for NP beyond the SM and analyze its contribution to the CP asymmetry of $B \to K\pi$ in SCET. As mentioned, the impact of SUSY appears only in the Wilson coefficients at the electroweak scale. Here we focus on the relevant contributions that may play an important role in the CP asymmetry of $B \to K\pi$, in particular, the gluino contribution to the chromomagnetic and EW penguins, namely $C^{\tilde g}_{8g}$, $C^{\tilde g}_7$, and $C^{\tilde g}_9$, and, in addition, the chargino contribution to the Z-penguin C^{χ}_9. These can be written in MIA as [22,23]:

$$C^{\tilde g}_{8g} \simeq \frac{8\alpha_S\pi}{9\sqrt{2}G_F V_{tb}V^*_{ts}m^2_{\tilde q}}\frac{m_{\tilde g}}{m_b}[(\delta^d_{LR})_{23} + (\delta^d_{RL})_{23}] \times \left(\frac{1}{3}M_1(x) + 3M_2(x)\right), \tag{39}$$

$$C^{\tilde g}_{7\gamma} \simeq \frac{\pi\alpha_S}{6\sqrt{2}G_F V_{tb}V^*_{ts}m^2_{\tilde q}}\frac{N^2_c - 1}{2N_c} \times \left[(\delta^d_{LL})_{23}\frac{1}{4}P_{1,3,2}(x,x) + (\delta^d_{RL})_{23}\frac{m_{\tilde g}}{m_b}P_{1,2,2}(x,x)\right], \tag{40}$$

$$C^{\tilde g}_9 \simeq -\frac{\pi\alpha_S}{6\sqrt{2}G_F V_{tb}V^*_{ts}m^2_{\tilde q}}\frac{N^2_c - 1}{2N_c}(\delta^d_{LL})_{23}\frac{1}{3}P_{0,4,2}(x,x), \tag{41}$$

$$C^{\chi}_9 \simeq \frac{\alpha}{4\pi}Y_t[(\delta^u_{RL})_{32} + \lambda(\delta^u_{RL})_{31}] \times \left(4\left(1 - \frac{1}{4\sin^2\theta_W}\right)R_C + R_D\right), \tag{42}$$

where $x = m^2_{\tilde g}/m^2_{\tilde q}$ and the functions $M_{1,2}$, P_{ijk} and $R_{C,D}$ are the corresponding loop functions, which depend on SUSY parameters through gluino/chargino mass and squark masses and can be found in Refs. [22,23]. Note that although $(\delta^d_{LR})_{23}$ and $(\delta^d_{RL})_{23}$ are constrained by the experimental limits of $b \to s\gamma$ to be less than $\mathcal{O}(10^{-2})$, their contributions to $C^{\tilde g}_{8g}$ and $C^{\tilde g}_7$ are enhanced by a large factor of $m_{\tilde g}/m_b$. On the other hand, the mass insertion $(\delta^u_{RL})_{32}$ is free from any stringent constraints, and it can be of order one.

As advocated in the Introduction, SUSY models include new CP violating phases beyond the SM phase $\delta_{\rm CKM}$. These phases arise from the complex soft SUSY breaking terms. In MIA, the SUSY CP violating phases lead to complex mass insertions $(\delta^{u,d}_{AB})_{ij}$, hence complex SUSY Wilson coefficients, unlike in SM. A SUSY model with nonuniversal A-terms, which can be obtained in most of SUSY breaking scenarios, is the natural framework for inducing new SUSY sources of CP and flavor violation that yield observable effects in the low-energy CP violation experiments without exceeding the experimental EDM limits [12]. For $m_{\tilde g} = 300$ GeV and $m_{\tilde q} = 500$ GeV, the SUSY contributions to QCD and EW penguins can be approximated by

$$(\hat P)_{\rm SUSY} = (-0.004 + 0.0002i)(\delta^d_{LL})_{23} - 0.36(\delta^d_{LR})_{23} - 0.36(\delta^d_{RL})_{23} - 0.00004(\delta^u_{RL})_{32}, \tag{43}$$

$$(\hat{P}_{\rm EW})_{\rm SUSY} = (0.025 - 0.0005i)(\delta^d_{LL})_{23} + 0.00031(\delta^u_{RL})_{32}, \tag{44}$$

$$(\hat{P}^C_{\rm EW})_{\rm SUSY} = (0.013 - 0.0005i)(\delta^d_{LL})_{23} + 0.00013(\delta^u_{LR})_{32}. \tag{45}$$

Recall that the SM contribution to these parameters is given by

$$(\hat{P})_{\rm SM} = -0.006 + 0.0016i, \tag{46}$$

$$(\hat{P}_{\rm EW})_{\rm SM} = -0.0005 + 0.0001i, \tag{47}$$

$$(\hat{P}^C_{\rm EW})_{\rm SM} = -0.0002 + 0.0001i. \tag{48}$$

From the $b \to s\gamma$ constraints, one can fix the relevant mass insertions as follows:

$$(\delta^d_{LL})_{23} = e^{i\alpha^d_1}, \qquad (\delta^d_{LR})_{23} = (\delta^d_{RL})_{23} = 0.01 e^{i\alpha^d_2}, \qquad (\delta^u_{RL})_{32} = 1 e^{i\alpha^u}, \tag{49}$$

with unconstrained CP violating phases: $\alpha^d_{1,2}$ and α^u. It is clear that the QCD penguin is dominated by the SM contribution, which is essentially the charm penguin effect. However, the EW penguins, which are quite suppressed in the SM, receive significant contributions in the SUSY models, in particular, due to the gluino contribution to the EW penguin with photon mediation. In this case, one can approximate $r_{\rm EW}$ and $r^C_{\rm EW}$ as

$$r_{\rm EW} = (r_{\rm EW})^{\rm SM} \left| \frac{1 - 46.5 e^{i\alpha^d_1} - 0.58 e^{i\alpha^u}}{1 + (0.65 + 0.13i) e^{i\alpha^d_1} + (1.64 + 0.41i) e^{i\alpha^d_2}} \right|, \tag{50}$$

$$r^C_{\rm EW} \simeq (r^C_{\rm EW})^{\rm SM} \left| \frac{1 - 50.7 e^{i\alpha^d_1} - 0.52 e^{i\alpha^u}}{1 + (0.65 + 0.13i) e^{i\alpha^d_1} + (1.64 + 0.41i) e^{i\alpha^d_2}} \right|. \tag{51}$$

From these expressions, it is clear that the magnitudes of $r_{\rm EW}$ and $r^C_{\rm EW}$ can be significantly enhanced and reach up to tens of the SM results. As we concluded in the previous section, a large value of $r_{\rm EW}$ and/or $r^C_{\rm EW}$, besides nonvanishing CP violating phases $\theta_{\rm EW}$ and $\theta^C_{\rm EW}$, is an essential condition for resolving the $B \to K\pi$ puzzle. Also one notes that the chargino exchange gives subdominant contribution.

One can also notice that the relation $r_{\rm EW} \sim 2r^C_{\rm EW}$ remains valid in SUSY models, as in the SM. Furthermore, since the mass insertion $(\delta^d_{LL})_{23}$ gives the dominant contributions to $P_{\rm EW}$ and $P^C_{\rm EW}$, one gets $\sin(\theta_P - \theta_{\rm EW}) \sim \sin(\theta_P - \theta^C_{\rm EW})$. Therefore, the condition of accounting for the discrepancy in $B \to K\pi$ CP asymmetries, Eq. (38), leads to

$$r_{\rm EW} \sin(\theta_P - \theta_{\rm EW}) \left(1 - \frac{2r^C_{\rm EW}}{r_{\rm EW}}\right) \simeq \frac{0.12}{\sin\delta_P} \sim 0.4, \tag{52}$$

where $\sin(\theta_P - \theta_{\rm EW}) \sim \mathcal{O}(1)$ and $(1 - 2r_{\rm EW}/r^C_{\rm EW}) \sim \mathcal{O}(0.1)$. Therefore, the CP asymmetries of $B \to K\pi$ can be accommodated if $r_{\rm EW} \geq \mathcal{O}(1)$, which can be obtained as shown in Eq. (50).

As an example, one can check that the following values of the mass insertion phases, $\alpha^d_1 = 2.1$ rad, $\alpha^d_2 = 1.5$ rad, and $\alpha^u = 0$, lead to $r_{\rm EW} \simeq 1.7$ and $r^C_{\rm EW} = 0.9$. This means that both $r_{\rm EW}$ $r^C_{\rm EW}$ are enhanced from the SM result by a factor of 20. Also, in this case, one finds the SUSY CP violating phases as follows: $\theta_{\rm EW} = -2.25$ rad and $\theta^C_{\rm EW} = -2.27$ rad. These results imply that the CP asymmetries of $B \to K^+\pi^0$ and $B \to K^+\pi^-$ are given by

$$A^{CP}_{K^+\pi^0} = 0.06, \qquad A^{CP}_{K^+\pi^-} = -0.09, \tag{53}$$

which are in agreement with the experimental measurements reported in Table I. It is important to note that since $r_{\rm EW} \ll 1$, one must use the complete expression for the CP asymmetries to get the correct results.

VI. CONCLUSIONS

In this work, we have studied the large discrepancy in the experimentally measured asymmetries of $B \to K\pi$ in the SCET framework. We conclude that in the standard model, one cannot accommodate all the experimental results in the SCET framework.

We have considered the possibility that New Physics could satisfy the measured asymmetries. We have classified the properties of New Physics needed to bring the theoretical results to an experimentally acceptable level in the SCET scenario. A general feature is that a new source of CP violation must emerge. As an example of a New Physics model, we studied supersymmetric models with minimal particle content in a model independent fashion by utilizing mass insertion approximation. We found that the gluino contribution to the electroweak penguin is essential. In our analysis we let trilinear A-terms vary freely, in which case we can find an experimentally allowed region in the parameter space.

Therefore, if SCET is a reliable way to treat hadronic matrix elements, the present experimental results indicate New Physics. Supersymmetric models remain a viable candidate for such New Physics, if nonminimal flavor violation is allowed.

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