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Observation of CP violation in $D^0 \to K^-\pi^+$ as a smoking gun for new physics

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In this paper, we study the Cabibbo favored nonleptonic D^0 decays into $K^-\pi^+$ decays. First we show that, within the Standard Model (SM), the corresponding charge conjugation and parity (CP) asymmetry is strongly suppressed and out of the experimental range, even taking into account the large strong phases coming from final state interactions. We show also that although new physics models with extra sequential generation can enhance the CP asymmetry by a few orders of magnitude, the resulting CP asymmetry is still far from the experimental range. The most sensitive new physics models to this CP asymmetry come from nonmanifest left-right models, where a CP asymmetry up to 10% can be reached, and the general two-Higgs model extension of the SM, where a CP asymmetry of order 10^{-2} can be obtained without being in contradiction with the experimental constraints on these models.

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I. INTRODUCTION

The Standard Model (SM) has been very successful in predicting and fitting all the experimental measurements up to date over energies ranging many orders of magnitude [1]. Unfortunately, the SM is only a patchwork where several sectors remain totally unconnected. Flavor physics, for example, involves quark masses, mixing angles, and charge conjugation and parity (CP) violating phases appearing in the Cabibbo-Kobayashi-Maskawa (CKM) quark mixing matrix [2,3]. These parameters unavoidably have to be measured and are independent from parameters present in other sectors like electroweak symmetry breaking, quantum chromodynamics, etc. Other sectors remain to be tested like CP violation in the up quark sector and even tensions with experimental measurements remain to be clarified (see for instance Refs. [4–7]).

This is why it is important to find processes where the SM predictions are very well known and a simple measurement can show their discrepancy. Some of these processes are the rare decays and other "null" tests which correspond to an observable strictly equal to zero within the SM. So any deviation from zero of these null test observables is a clear signal of physics beyond the SM. This is the case for Cabibbo-favored (CF) and double Cabibbo suppressed (DCS) nonleptonic charm decays, where the direct CP violation is very suppressed given that penguin diagrams are absent [8–10].

Even with the observation of D^0 oscillation [11–16] and the first signal of CP violation in $D \to 2\pi$, $2K$ [singly Cabibbo suppressed (SCS) modes] [17–32], it is not clear that the SM [33–41] can describe correctly the CP violation in the up quark sector. It is even more difficult as large distance contributions are important and difficult to be evaluated [42–46]. Thus, unfortunately it will be not easy to find new physics in this sector (SCS). On the contrary, in the CF and Double Cabibbo Favored modes, the charge conjugation and parity violation (CPV) signal is very suppressed in the SM and in most of the new physics (NP) models. So even if long distance corrections are large, any small signal of CPV will be due to new physics.

Up to now, only $D^0 \leftrightarrow \bar{D}^0$ oscillations have been observed and their parameters have been measured [1,11–16]:

$$x \equiv \frac{\Delta m_d}{\Gamma_D} = 0.55^{+0.12}_{-0.13}, \qquad y \equiv \frac{\Delta\Gamma_D}{2\Gamma_D} = 0.83(13), \quad (1)$$

$$\left|\frac{q}{p}\right| = 0.91^{0.18}_{0.16}, \quad \phi \equiv \arg(q/p) = -(10.2^{+9.4}_{-8.9})^\circ, \quad (2)$$

where $x \neq 0$ and/or $y \neq 0$ mean oscillations have been observed, while $|q/p| \neq 1$ and/or $\phi \neq 0$ are necessary to have CP violation. The theoretical estimations of these parameters [1] are not easy, as they have large uncertainties given that the c quark is not heavy enough to apply heavy quark effective (HQE) theory (like in B physics) [47]. Similarly, it is not light enough to use chiral perturbation theory (CPTh) (like in kaon physics). Besides, there are cancellations due to the Glashow-Iliopoulos-Maiani mechanism [2,48]. Theoretically CP violation in the charm sector is smaller than in the B and kaon sectors. This is due to a combination of factors: CKM matrix elements ($|V_{ub}V^*_{cb}/V_{us}V^*_{cs}|^2 \sim 10^{-6}$) and the fact that the b quark mass is small compared to the top mass. CP violation in the b quark sector is due to the large top quark mass, while in the kaon it is due to a combination of the charm and top quark.

Experimental data should be improved within the next years with LHCb [49] and the different charm factory

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projects [50]. In Table I the experimentally measured branching ratios and CP asymmetries are given for different nonleptonic D decays.

In this paper, we study in detail the CP asymmetry for the CF $D^0 \to K^-\pi^+$ decay. In Sec. II, we give the general description of the effective Hamiltonian describing this decay within the SM and show how to evaluate the strong phases needed to get CP violating observables. These strong phases are generated through final state interaction (FSI). In Sec. III, we first evaluate the SM prediction for the CP asymmetry and we show that within SM, such CP asymmetry is experimentally out of range. In Sec. IV, new physics models are introduced and their contributions to CP asymmetry are evaluated. Finally, we conclude in Sec. V.

II. GENERAL DESCRIPTION OF CF NONLEPTONIC D^0 DECAYS INTO K^- AND π^+

In general the Hamiltonian describing $D^0 \to K^-\pi^+$ is given by

$$\mathcal{L}_{\text{eff.}} = \frac{G_F}{\sqrt{2}} V_{cs}^* V_{ud}\left[\sum_{i,a} c^i_{1ab}\bar{s}\Gamma^i c_a \bar{u}\Gamma_i d_b + \sum_{i,a} c^i_{2ab}\bar{u}\Gamma^i c_a \bar{s}\Gamma_i d_b\right], \quad (3)$$

with $i = $ S, V, and T for the scalar (S), vectorial (V), and tensorial (T) operators, respectively. The Latin indexes a, $b = L, R$ and $q_{L,R} = (1 \mp \gamma_5)q$.

Within the SM, only two operators contribute to the effective Hamiltonian for this process [8–10]. The other operators can only be generated through new physics.

$$\mathcal{H} = \frac{G_F}{\sqrt{2}} V_{cs}^* V_{ud}(c_1 \bar{s}\gamma_\mu c_L \bar{u}\gamma^\mu d_L + c_2 \bar{u}\gamma_\mu c_L \bar{s}\gamma^\mu d_L) + \text{H.c.} \quad (4)$$

$$= \frac{G_F}{\sqrt{2}} V_{cs}^* V_{ud}(c_1\mathcal{O}_1 + c_2\mathcal{O}_2) + \text{H.c.}, \quad (5)$$

where $a_1 \equiv c_1 + c_2/N_c = 1.2 \pm 0.1$ and $a_2 \equiv c_2 - c_1/N_C = -0.5 \pm 0.1$ [8–10], where N_C is the color number. For the case $D \to K\pi$ [8–10], one has that

$$A_{D^0\to K^-\pi^+} = -i\frac{G_F}{\sqrt{2}} V_{cs}^* V_{ud}[a_1 X^{\pi^+}_{D^0K^-} + a_2 X^{D^0}_{K^-\pi^+}], \quad (6)$$

$$\text{BR} = \frac{\tau_D p_K}{8\pi m_D^2}|A|^2, \quad (7)$$

where BR is the branching ratio of the process. τ_D is the D lifetime, p_K is the kaon momentum and m_D is the D meson mass. The $X^{\pi^+}_{D^0K^-}$ and $X^{D^0}_{K^-\pi^+}$ can be expressed in the following way:

$$X^{P_1}_{P_2P_3} = if_{P_1}\Delta^2_{P_2P_3}F_0^{P_2P_3}(m^2_{P_1}), \qquad \Delta^2_{P_2P_3} = m^2_{P_2} - m^2_{P_3}, \quad (8)$$

where f_D and f_K are the decay constants for D and K mesons, respectively, and F_0^{DK} and $F_0^{D\pi}$ are the corresponding form factors. These amplitudes have been computed within the so-called naive factorization approximation (NFA) without including the FSI. In NFA, no strong CP conserving phases are obtained (and therefore no CPV is predicted) but it is well known that FSI effects are very important in these channels [52–56]. In principle, we have many FSI contributions: resonances, other intermediate states, rescattering, and so on. Resonances are especially important in this region given that they are abundant. They can be included and seem to produce appropriate strong phases [56]. However, the other contributions mentioned above have to be included too, rendering the theoretical prediction cumbersome. A more practical approach, although less predictive, is obtained by fitting the experimental data [52,56]. This is the so-called quark diagram

TABLE I. Direct CP in D nonleptonic decays, from the Heavy Flavor Averaging Group (HFAG) [1,51]. The blank entries correspond to cases where no experimental data for branching ratio or CP asymmetries are available at present time.

Mode	BR (%)	A_{CP} (%)	Mode	BR (%)	A_{CP} (%)
$D^0 \to K^-\pi^+$ CF	3.95(5)		$D^0 \to \bar{K}^0\pi^0$ CF	2.4(1)	
$D^0 \to \bar{K}^0\eta$ CF	0.96(6)		$D^0 \to \bar{K}^0\eta'$ CF	1.90(11)	
$D^+ \to \bar{K}^0\pi^+$ CF	3.07(10)		$D_s^+ \to K^+\bar{K}^0$ CF	2.98(8)	
$D_s^+ \to \pi^+\eta$ CF	1.84(15)		$D_s^+ \to \pi^+\eta'$ CF	3.95(34)	
$D^0 \to K^+\pi^-$ DCS	$1.48(7) \times 10^{-4}$		$D^0 \to K^0\pi^0$ DCS		
$D^0 \to K^0\eta$ DCS			$D^0 \to K^0\eta'$ DCS		
$D^+ \to K^0\pi^+$ DCS			$D^+ \to K^+\pi^0$ DCS	$1.72(19) \times 10^{-2}$	
$D^+ \to K^+\eta$ DCS	$1.08(17) \times 10^{-2}$		$D^+ \to K^+\eta'$ DCS	$1.76(22) \times 10^{-2}$	
$D_s^+ \to K^+K^0$ DCS					
$D^0 \to \pi^-\pi^+$	0.143(3)	0.22(24)(11)			
$D^0 \to K^-K^+$	0.398(7)	−0.24(22)(9)	$A_{CP}(K^+K^-) - A_{CP}(\pi^+\pi^-)$		−0.65(18)
$D^+ \to K_S^0\pi^+$	1.47(7)	−0.71(19)(20)	$D^\pm \to \pi^+\pi^-\pi^\pm$	0.327(22)	1.7(42)
$D^\pm \to K^\mp\pi^\pm\pi^\pm$	9.51(34)	−0.5(4)(9)	$D^\pm \to K_s^0\pi^\pm\pi^0$	6.90(32)	0.3(9)(3)
$D^\pm \to K^+K^-\pi^\pm$	0.98(4)	0.39(61)			

approach. Within this approach, the amplitude is decomposed into parts corresponding to generic quark diagrams. The main contributions are the tree-level quark contribution (T) and exchange quark diagrams (E). Their results can be summarized in the following way, for the process under consideration [56]:

$$A_{D^0\to K^-\pi^+} \equiv V_{cs}^* V_{ud}(T+E), \tag{9}$$

with

$$\begin{aligned} T &= (3.14 \pm 0.06) \times 10^{-6}\ \mathrm{GeV} \\ E &= 1.53^{+0.07}_{-0.08} \times 10^{-6} \cdot \mathrm{e}^{(122\pm 2)^\circ i}\ \mathrm{GeV}, \end{aligned} \tag{10}$$

where in NFA they can be approximately written as

$$T \simeq \frac{G_F}{\sqrt{2}} a_1 f_\pi (m_D^2 - m_K^2) F_0^{DK}(m_\pi^2) \tag{11}$$

$$E \simeq -\frac{G_F}{\sqrt{2}} a_2 f_D (m_K^2 - m_\pi^2) F_0^{K\pi}(m_D^2). \tag{12}$$

In the rest of this work we are going to use the values obtained by the experimental fit, given in Eq. (10).

III. *CP* ASYMMETRY IN $D^0 \to K^-\pi^+$ WITHIN SM

In the case of CF (and DCS) processes the corrections are very small (see diagrams in Figs. 1 and 2) and are generated through box and dipenguin diagrams [57–59]. In this section, we shall evaluate these contributions.

The box contribution is given as [59,60]

$$\Delta\mathcal{H} = \frac{G_F^2 m_W^2}{2\pi^2} V_{cD}^* V_{uD} V_{Us}^* V_{Ud} f(x_U, x_D) \bar{u}\gamma_\mu c_L \bar{s}\gamma^\mu d_L \tag{13}$$

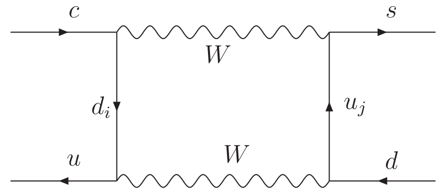

FIG. 1. Feynman diagram for CF processes: box contribution.

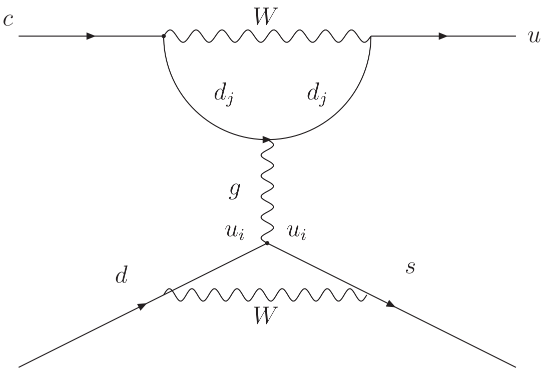

FIG. 2. Feynman diagram for CF processes: dipenguin contribution.

$$= \frac{G_F^2 m_W^2}{2\pi^2} \lambda_{cu}^D \lambda_{sd}^U f(x_U, x_D) \mathcal{O}_2 = \frac{G_F^2 m_W^2}{2\pi^2} b_x \mathcal{O}_2, \tag{14}$$

where

$$b_x \equiv \lambda_{cu}^D \lambda_{sd}^U f(x_U, x_D) \tag{15}$$

$$\begin{aligned} &= V_{cd}^* V_{ud}(V_{us}^* V_{ud} f_{ud} + V_{cs}^* V_{cd} f_{cd} + V_{ts}^* V_{td} f_{td}) + V_{cs}^* V_{us}(V_{us}^* V_{ud} f_{us} + V_{cs}^* V_{cd} f_{cs} + V_{ts}^* V_{td} f_{ts}) \\ &\quad + V_{cb}^* V_{ub}(V_{us}^* V_{ud} f_{ub} + V_{cs}^* V_{cd} f_{cb} + V_{ts}^* V_{td} f_{tb}) \\ &= V_{cs}^* V_{us}[V_{cs}^* V_{cd}(f_{cs} - f_{cd} - f_{us} + f_{ud}) + V_{ts}^* V_{td}(f_{ts} - f_{td} - f_{us} + f_{ud})] \\ &\quad + V_{cb}^* V_{ub}[V_{cs}^* V_{cd}(f_{cb} - f_{cd} - f_{ub} + f_{ud}) + V_{ts}^* V_{td}(f_{tb} - f_{td} - f_{ub} + f_{ud})], \end{aligned} \tag{16}$$

with $\lambda_{DD'}^U \equiv V_{UD}^* V_{UD'}$, $\lambda_{UU'}^D \equiv V_{UD}^* V_{U'D}$, $U = u, c, t$, $D = d, s, b$, $x_q = (m_q/m_W)^2$, $f_{UD} \equiv f(x_U, x_D)$ [61], and

$$f(x,y) = \frac{7xy - 4}{4(1-x)(1-y)} + \frac{1}{x-y}\left[\frac{y^2 \log y}{(1-y)^2}\left(1 - 2x + \frac{xy}{4}\right) - \frac{x^2 \log x}{(1-x)^2}\left(1 - 2y + \frac{xy}{4}\right)\right].$$

Numerically, one obtains

$$b_x \simeq 3.6 \times 10^{-7} \mathrm{e}^{0.07\cdot i}. \tag{17}$$

The quark masses' values are taken at the m_c scale as given in Ref. [1]. The other contribution to the Lagrangian is the dipenguin and it gives [57,58,62]

$$\begin{aligned} \Delta\mathcal{H} &= -\frac{G_F^2 \alpha_S}{8\pi^3}[\lambda_{cu}^D E_0(x_D)][\lambda_{sd}^U E_0(x_U)]\bar{s}\gamma_\mu T^a d_L (g^{\mu\nu}\Box - \partial^\mu\partial^\nu)\bar{u}\gamma_\nu T^a c_L \\ &= -\frac{G_F^2 \alpha_S}{8\pi^3} p g \bar{s}\gamma_\mu T^a d_L (g^{\mu\nu}\Box - \partial^\mu\partial^\nu)\bar{u}\gamma_\nu T^a c_L \equiv \frac{G_F^2 \alpha_S}{16\pi^3} p_g \mathcal{O} \end{aligned} \tag{18}$$

$$\begin{aligned} p_g &\equiv [\lambda^D_{cu} E_0(x_D)][\lambda^U_{sd} E_0(x_U)] \\ &= [V^*_{cs} V_{us}(E_0(x_s) - E_0(x_d)) \\ &\quad + V^*_{cb} V_{ub}(E_0(x_b) - E_0(x_d))] \\ &\quad \times [V_{cd} V^*_{cs}(E_0(x_c) - E_0(x_u)) \\ &\quad + V_{td} V^*_{ts}(E_0(x_t) - E_0(x_u))], \end{aligned} \tag{19}$$

where T^a are the generators of $SU(3)_C$. Numerically, $p_g \simeq -1.62 \cdot e^{-0.002i}$ and the Inami functions are given by

$$\begin{aligned} E_0(x) = \frac{1}{12(1-x)^4}[x(1-x)(18-11x-x^2) \\ - 2(4-16x+9x^2)\log(x)]. \end{aligned} \tag{20}$$

The operator $\mathcal{O}$ can be reduced as

$$\begin{aligned} \mathcal{O} &= \bar{s}\gamma_\mu T^a d_L(g^{\mu\nu}\Box - \partial^\mu\partial^\nu)\bar{u}\gamma_\nu T^a c_L = \bar{s}\gamma_\mu T^a d_L \Box(\bar{u}\gamma^\nu T^a c_L) + \bar{s}\not{\partial} T^a d_L \bar{u}\not{\partial} T^a c_L \\ &= -q^2 \bar{s}\gamma_\mu T^a d_L \bar{u}\gamma^\mu T^a c_L - (m_s \bar{s} T^a d_{S-P} + m_d \bar{s} T^a d_{S+P}) \cdot (m_c \bar{u} T^a c_{S+P} + m_u \bar{u} T^a c_{S-P}) - q^2 \bar{s}\gamma_\mu T^a d_L \bar{u}\gamma^\mu T^a c_L \\ &\quad - m_s m_c \bar{s} T^a d_L \bar{u} T^a c_R - m_d m_u \bar{s} T^a d_R \bar{u} T^a c_L - m_s m_u \bar{s} T^a d_L \bar{u} T^a c_L - m_d m_c \bar{s} T^a d_R \bar{u} T^a c_R, \end{aligned} \tag{21}$$

where q^2 is the gluon momentum and N is the color number. This expression can be simplified using the fact that

$$\begin{aligned} \bar{s}\gamma_\mu T^a d_L \bar{u}\gamma^\mu T^a c_L &= \frac{1}{2}\left(\mathcal{O}_1 - \frac{1}{N}\mathcal{O}_2\right) \\ \bar{s} T^a d_L \bar{u} T^a c_R &= -\frac{1}{4}\bar{s}\gamma_\mu c_R \bar{u}\gamma^\mu d_L - \frac{1}{2N}\bar{s} d_L \bar{u} c_R \\ \bar{s} T^a d_R \bar{u} T^a c_L &= -\frac{1}{4}\bar{s}\gamma_\mu c_L \bar{u}\gamma^\mu d_R - \frac{1}{2N}\bar{s} d_R \bar{u} c_L \\ \bar{s} T^a d_L \bar{u} T^a c_L &= -\frac{1}{4}\bar{s} c_L \bar{u} d_L - \frac{1}{16}\bar{s}\sigma_{\mu\nu} c_L \bar{u}\sigma^{\mu\nu} d_L - \frac{1}{2N}\bar{s} d_L \bar{u} c_L \\ \bar{s} T^a d_R \bar{u} T^a c_R &= -\frac{1}{4}\bar{s} c_R \bar{u} d_R - \frac{1}{16}\bar{s}\sigma_{\mu\nu} c_R \bar{u}\sigma^{\mu\nu} d_R - \frac{1}{2N}\bar{s} d_R \bar{u} c_R. \end{aligned} \tag{22}$$

Once taking the expectation values, one obtains

$$\begin{aligned} \langle\mathcal{O}\rangle &= -q^2\langle\bar{s}\gamma_\mu T^a d_L \bar{u}\gamma^\mu T^a c_L\rangle - m_s m_c\langle\bar{s} T^a d_L \bar{u} T^a c_R\rangle - m_d m_u\langle\bar{s} T^a d_R \bar{u} T^a c_L\rangle - m_s m_u\langle\bar{s} T^a d_L \bar{u} T^a c_L\rangle \\ &\quad - m_d m_c\langle\bar{s} T^a d_R \bar{u} T^a c_R\rangle \\ &\simeq -\frac{q^2}{2}\left(1 - \frac{1}{N^2}\right) X^{\pi^+}_{D^0 K^-} + \frac{m_s m_c}{4}\left(1 - \frac{1}{N}\right) X^{\pi^+}_{D^0 K^-} + \frac{5 m_d}{8 N m_s} m_D^2 X^{D^0}_{K^-\pi^+} \end{aligned} \tag{23}$$

Hence, one gets for the Wilson coefficients

$$\begin{aligned} \Delta a_1 &= -\frac{G_F m_W^2}{\sqrt{2}\pi^2 V^*_{cs} V_{ud} N} b_x - \frac{G_F \alpha_S}{4\sqrt{2}\pi^3 V_{cs} V^*_{us}}\left[\frac{q^2}{2}\left(1 - \frac{1}{N^2}\right) - \frac{m_c m_s}{4}\left(1 - \frac{1}{N}\right)\right] p_g \simeq 2.8 \times 10^{-8} e^{-0.004i} \\ \Delta a_2 &= -\frac{G_F m_W^2}{\sqrt{2}\pi^2 V^*_{cs} V_{ud}} b_x - \frac{G_F \alpha_S}{4\sqrt{2}\pi^3 V_{cs} V^*_{us}} \frac{5 m_d m_D^2}{8 N m_s} p_g \simeq -2.0 \times 10^{-9} e^{0.07i}, \end{aligned} \tag{24}$$

where to obtain the last result we have used the fact that for the decay $D^0 \to K^-\pi^+$, one can approximate $q^2 = (p_c \mp p_u)^2 = (p_s \pm p_d)^2 \simeq (p_D - p_\pi/2)^2 = (m_D^2 + m_K^2)/2 + 3m_\pi^2/4$, by assuming that $p_c \simeq p_D$, $p_u \simeq p_\pi/2$, and $\alpha_S \simeq 0.3$. It should be noticed that the box contribution is dominated by the heavy quarks while the penguin is dominated by the light ones. The direct CP asymmetry is then

$$\begin{aligned} A_{CP} &= \frac{|A|^2 - |\bar{A}|^2}{|A|^2 + |\bar{A}|^2} = \frac{2|r|\sin(\phi_2 - \phi_1)\sin(\alpha_E)))}{|1 + r|^2} \\ &= 1.4 \times 10^{-10}, \end{aligned} \tag{25}$$

with $r = E/T$, $a_i \to a_i + \Delta a_i = a_i + |\Delta a_i| \exp[i\Delta\phi_i]$, and $\phi_i \simeq \Delta a_i \sin\Delta\phi_i / a_i$, and α_E is the conserving phase which appears in Eq. (10).

IV. NEW PHYSICS

With new physics, the general Hamiltonian is not only given by $\mathcal{O}_{1,2}$. The expressions of the expectation values of these operators can be found in the Appendix. It is important to notice that as expected, only two form factors appear, namely, $\chi^{D^0}_{K^-\pi^+}$ and $\chi^{\pi^+}_{D^0 K^-}$. Thus, it is important to take into account the FSI interactions, as the first one is identified as the E contribution and the second one is identified as the T contribution. In the next subsections, we shall calculate the Wilson coefficient for different models of new physics. For the first case, we will be assuming an extra SM fermion family. The second example will be to compute the CP asymmetry generated by a new charged gauge boson as it appears, for instance, in models based on gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, and our last subsection is dedicated to the effects of *CP* asymmetry coming from new charged Higgs-like scalar fields, applying it to the two-Higgs extension of the SM (type II and type III).

A. Contributions to A_{CP} from extra SM fermion family

A simple extension of the SM is the introduction of a new sequential generation of quarks and leptons (SM4). A fourth generation is not excluded by precision data [63–70]. Recent reviews on consequences of a fourth generation can be found in Refs. [71–81].

The $B \to K\pi$ *CP* asymmetry puzzle is easily solved by a fourth generation [82–84] with a mass within the following range [82]:

$$400\ \text{GeV} < m_{u_4} < 600\ \text{GeV}. \tag{26}$$

The values of SM4 parameters compatible with the high precision LEP measurements [64–66,69] are

$$m_{u_4} - m_{d_4} \simeq \left(1 + \frac{1}{5}\ln\frac{m_H}{115\ \text{GeV}}\right) \times 50\ \text{GeV} \tag{27}$$

$$|V_{ud_4}|, |V_{u_4 d}| \lesssim 0.04, \tag{28}$$

where V is the CKM quark mixing matrix which is now a 4×4 unitary matrix. The direct search limits from LEPII and CDF [85–87] are given by

$$m_{u_4} > 311\ \text{GeV} \qquad m_{d_4} > 338\ \text{GeV}. \tag{29}$$

Direct searches by the ATLAS and CMS collaborations have excluded $m_{d_4} < 480\,\text{GeV}$ and $m_{q_4} < 350\,\text{GeV}$ [88–90]. Thanks to the LHC, these limits are moving very quickly. Recently, ATLAS reported a new limit on $m_{u_4} > 656$ GeV at 95% confidence level [91], above the tree-level unitarity limit, $m_{u_4} < \sqrt{4\pi/3}v \simeq 504$ GeV. But SM4 is far from being completely understood. Most of the experimental constraints are model dependent. For instance, it has been shown in Ref. [92] that the bound on m_{u_4} should be relaxed up to $m_{u_4} > 350$ GeV if the decay $u_4 \to ht$ dominates. The recent LHC results which observe an excess in the $H \to \gamma\gamma$ corresponding to a Higgs mass around 125 GeV [93,94] seem to exclude the SM4 scenario [95,96], but these results are based on the fact that once we include the next-to-leading order electroweak corrections, the rate $\sigma(gg\to H)\times \text{Br}(H\to\gamma\gamma)$ is suppressed by more than 50% compared to the rate including only the leading order corrections [95,97–101]. This could be a signal of a nonperturbative regime which in SM4 can be easily reached at this scale due to the fourth generation strong Yukawa couplings. Therefore, direct and model-independent searches for fourth generation families at collider physics are still necessary to completely exclude the SM4 scenario.

The *CP* asymmetry in the model with a fourth family is easy to compute, as the contributions come from the same diagrams in the SM after adding just an extra $u_4 \equiv t'$ and $d_4 \equiv b'$. Similarly, in Ref. [90] it has been found that new CKM matrix elements can be obtained (all consistent with zero and for $m_{b'} = 600$ GeV) to be

$$\begin{gathered} s_{14} = |V_{ub'}| = 0.017(14), \qquad s_{24} = \frac{|V_{cb'}|}{c_{14}} = \frac{0.0084(62)}{c_{14}}, \qquad s_{34} = \frac{|V_{tb'}|}{c_{14}c_{24}} = \frac{0.07(8)}{c_{14}c_{24}} \\ |V_{t'd}| = |V_{t's}| = 0.01(1), \qquad |V_{t'b}| = 0.07(8), \qquad |V_{t'b'}| = 0.998(6), \qquad |V_{tb}| \geq 0.98 \qquad \tan\theta_{12} = \left|\frac{V_{us}}{V_{ud}}\right|, \\ s_{13} = \frac{|V_{ub}|}{c_{14}}, \qquad \delta_{13} = \gamma = 68° \qquad |V_{cb}| = |c_{13}c_{24}s_{23} - u_{13}^* u_{14} u_{24}^*| \simeq c_{13}c_{24}s_{23}. \end{gathered} \tag{30}$$

The two remaining phases (ϕ_{14} and ϕ_{24}) are unbounded. Thus, the absolute values of the CKM elements for the three families remain almost unchanged but not their phases. From these values one obtains

$$\begin{aligned} s_{13} &= 0.00415, \qquad s_{12} = 0.225, \qquad s_{23} = 0.04, \\ s_{14} &= 0.016, \qquad s_{24} = 0.006, \qquad s_{34} = 0.04. \end{aligned} \tag{31}$$

For a fourth sequential family, the maximal value for the *CP* violation is obtained as

$$A_{CP} \simeq -1.1 \times 10^{-7}, \tag{32}$$

where one uses $|V_{ub'}| = 0.06$, $|V_{cb'}| = 0.03$, $|V_{tb'}| = 0.25$, $\phi_{14} = -2.9$, $\phi_{24} = 1.3$.

This maximal value is obtained when the parameters mentioned above are varied in the range allowed by the experimental constraints, according to Eq. (30) in a "three sigma" range. The phases are varied in the whole range from $-\pi$ to π. Thus, one can obtain an enhancement of a thousand, which may be large but is still very far from the experimental possibilities.

B. A new charged gauge boson in left-right models

In this section, we shall look to see what could be the effect on the *CP* asymmetry coming from a new charged gauge boson coupled to quarks and leptons. As an example of such models, we apply our formalism to a well-known extension of the Standard Model based on extending the SM gauge group to include a gauge $SU(2)_R$ [102–106]. So now, our gauge group defining the electroweak interaction is given by $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. This SM extension has been extensively studied in previous works (see for instance Refs. [107–111]) and their parameters have been strongly constrained by experiments [1,112–116]. Recently, CMS [117,118] and ATLAS [119,120] at the LHC have improved the bound on the scale of the W_R gauge boson mass [121]. The new diagrams contributing to $D \to K\pi$ are similar to the SM tree-level diagrams where W_L is replaced by a W_R. These diagrams contribute to the effective Hamiltonian in the following way, assuming no mixing between W_L and W_R gauge bosons:

$$\mathcal{H}_{\rm LR} = \frac{G_F}{\sqrt{2}}\left(\frac{g_R m_W}{g_L m_{W_R}}\right)^2 V_{Rcs}^* \times V_{Rud}(c_1' \bar{s}\gamma_\mu c_R \bar{u}\gamma^\mu d_R + c_2' \bar{u}\gamma_\mu c_R \bar{s}\gamma^\mu d_R) + \text{H.c.} = \frac{G_F}{\sqrt{2}} V_{cs}^* V_{ud}(c_1 \mathcal{O}_1 + c_2 \mathcal{O}_2) + \text{H.c.}, \tag{33}$$

where g_L and g_R are the gauge $SU(2)_L$ and $SU(2)_R$ couplings, respectively. m_W and m_{W_R} are the $SU(2)_L$ and $SU(2)_R$ charged gauge boson masses, respectively. V_R is the quark mixing matrix which appears in the right sector of the Lagrangian similar to the CKM quark mixing matrix. This new contribution can enhance the SM prediction for the CP asymmetry, but it is still suppressed due to the limit on M_{W_R} which has to be of order 2.3 TeV [121] in the case of nonmixing left-right models.

In Refs. [122,123] it has been shown that the mixing between the left and the right gauge bosons can strongly enhance any CP violation in the charm and muon sectors. This left-right (LR) mixing is restricted by deviation to nonunitarity of the CKM quark mixing matrix. The results were that the LR mixing angle called ξ has to be smaller than 0.005 [124] and the right scale M_R bigger than 2.5 TeV [121]. If the left-right symmetry is not manifest (essentially, that g_R could be different from g_L at the unification scale), the limit on the M_R scale is much less restrictive and the right gauge bosons could be as light as 0.3 TeV [125]. In such a case, ξ can be as large as 0.02, if large CP violation phases in the right sector are present [109], and still be compatible with experimental data [123,126,127]. Recently, precision measurement of the muon decay parameters done by the TWIST Collaboration [128,129] put a model-independent limit on ξ to be smaller than 0.03 (taking $g_L = g_R$). Let us now compute the effect of the LR mixing gauge boson on our CP asymmetry. So first, one defines the charged current mixing matrix [122]

$$\begin{pmatrix} W_L \\ W_R \end{pmatrix} = \begin{pmatrix} \cos\xi & -\sin\xi \\ e^{i\omega}\sin\xi & e^{i\omega}\cos\xi \end{pmatrix}\begin{pmatrix} W_1 \\ W_2 \end{pmatrix} \simeq \begin{pmatrix} 1 & -\xi \\ e^{i\omega}\xi & e^{i\omega} \end{pmatrix}\begin{pmatrix} W_1 \\ W_2 \end{pmatrix}, \tag{34}$$

where W_1 and W_2 are the mass eigenstates and $\xi \sim 10^{-2}$. Thus, the charged current interaction parts become

$$\mathcal{L} \simeq -\frac{1}{\sqrt{2}}\bar{U}\gamma_\mu(g_L V P_L + g_R \xi \bar{V}^R P_R) D W_1^\dagger - \frac{1}{\sqrt{2}}\bar{U}\gamma_\mu(-g_L \xi V P_L + g_R \bar{V}^R P_R) D W_2^\dagger, \tag{35}$$

where $V = V_{\rm CKM}$ and $\bar{V}^R = e^{i\omega} V^R$. Once one integrates out the W_1 in the usual way and, neglecting the W_2 contributions given its mass is much higher, one obtains the effective Hamiltonian responsible for our process:

$$\mathcal{H}_{\rm eff} = \frac{4G_F}{\sqrt{2}}\left[c_1 \bar{s}\gamma_\mu\left(V^* P_L + \frac{g_R}{g_L}\xi \bar{V}^{R*} P_R\right)_{cs} \times c\bar{u}\gamma^\mu\left(V P_L + \frac{g_R}{g_L}\xi \bar{V} P_R\right)_{ud} d + c_2 \bar{s}_\alpha \gamma_\mu\left(V^* P_L + \frac{g_R}{g_L}\xi \bar{V}^{R*} P_R\right)_{cs} \times c_\beta \bar{u}_\beta \gamma^\mu\left(V P_L + \frac{g_R}{g_L}\xi \bar{V} P_R\right)_{ud} d_\alpha\right] + \text{H.c.}, \tag{36}$$

where α, β are color indices. It is easy to check that taking the limit $\xi \to 0$, one obtains Eq. (5), with the only difference coming from the c_2 terms, because the Fierz transformation has been applied. The terms of the effective Hamiltonian proportional to ξ are

$$\Delta\mathcal{H}_{\rm eff} \simeq \frac{G_F}{\sqrt{2}}\frac{g_R}{g_L}\xi[c_1 \bar{s}\gamma_\mu V_{cs}^* c_L \bar{u}\gamma^\mu \bar{V}_{ud}^R d_R + c_1 \bar{s}\gamma_\mu \bar{V}_{cs}^{R*} c_R \bar{u}\gamma^\mu V_{ud} d_L + c_2 \bar{s}_\alpha \gamma_\mu V_{cs}^* c_{L\beta} \bar{u}_\beta \gamma^\mu \bar{V}_{ud}^R d_{R\alpha} + c_2 \bar{s}_\alpha \gamma_\mu \bar{V}_{cs}^{R*} c_{R\beta} \bar{u}_\beta \gamma^\mu V_{ud} d_{L\alpha}] + \text{H.c.}, \tag{37}$$

The contribution to the amplitude proportional to ξ is then given by

$$\Delta A = -\frac{iG_F}{\sqrt{2}}\frac{g_R}{g_L}\xi\Big[-c_1 V_{cs}^* \bar{V}_{ud}^R\Big(X_{D^0K^-}^{\pi^+} + \frac{2}{N}\chi^{D^0} X_{K^-\pi^+}^{D^0}\Big) + c_1 \bar{V}_{cs}^{R*} V_{ud}\Big(X_{D^0K^-}^{\pi^+} + \frac{2}{N}\chi^{D^0} X_{K^-\pi^+}^{D^0}\Big) - c_2 V_{cs}^* \bar{V}_{ud}^R\Big(2\chi^{D^0} X_{K^-\pi^+}^{D^0} + \frac{1}{N} X_{D^0K^-}^{\pi^+}\Big) + c_2 \bar{V}_{cs}^{R*} V_{ud}\Big(2\chi^{D^0} X_{K^-\pi^+}^{D^0} + \frac{1}{N} X_{D^0K^-}^{\pi^+}\Big)\Big] = \frac{iG_F}{\sqrt{2}}\frac{g_R}{g_L}\xi(V_{cs}^* \bar{V}_{ud}^R - \bar{V}_{cs}^{R*} V_{ud})(a_1 X_{D^0K^-}^{\pi^+} + 2\chi^{D^0} a_2 X_{K^-\pi^+}^{D^0}) = -\frac{g_R}{g_L}\xi(\bar{V}_{cs}^{R*} V_{ud} - V_{cs}^* \bar{V}_{ud}^R)(T - 2\chi^{D^0} E), \tag{38}$$

where χ^{π^+} and χ^{D^0} are defined as

$$\chi^{\pi^+} = \frac{m_\pi^2}{(m_c - m_s)(m_u + m_d)}$$
$$\chi^{D^0} = \frac{m_D^2}{(m_c + m_u)(m_s - m_d)}. \quad (39)$$

The *CP* asymmetry becomes

$$A_{CP} = \frac{4(g_R/g_L)\xi}{V_{cs}^* V_{ud}|1+r|^2}(1+2\chi^{D^0})\mathrm{Im}(\bar{V}_{cs}^{R*}V_{ud} - V_{cs}^*\bar{V}_{ud}^R)\mathrm{Im}(r), \quad (40)$$

with $r = E/T$. For a value as large as $\xi \sim 10^{-2}$, the asymmetry can be as large as 0.1. Also, we should note that to obtain these results, we have used the fact that the chiralities do not mix under strong interactions if the quark masses are not taken into account. This is approximately the case in the evolution of the Wilson coefficients from m_W to m_c, as the quarks in the loop are the down quarks; this is contrary to a process like $b \to s\gamma$, where the quarks in the QCD corrections are the up quarks and, in that case, a strong effect from top quarks could be expected [130–133]. In our case, as a first approximation, the QCD corrections to the Wilson coefficient coming from the running of the renormalization group from m_W to m_c can be safely neglected.

C. Models with charged Higgs contributions

Our last example of new physics is to consider the contribution to the effective Hamiltonian responsible for the $D^0 \to K^-\pi^+$ process due to new charged Higgs fields. The simple SM extensions which include new charged Higgs fields are the two-Higgs doublet models (2HDMs) [134,135]. Usually, it is used to classify these 2HDMs in three types: types I, II, or III (for a review, see Ref. [136]). In 2HDM type II models (like the minimal supersymmetric Standard Model), one Higgs couples to the down quarks and charged leptons and the other Higgs couples to up type quarks. LEP has performed a direct search for a charged Higgs in type II 2HDM and they obtained a bound of 78.6 GeV [137]. Recent results on $B \to \tau\nu$ obtained by BELLE [5] and *BABAR* [6] have strongly improved the indirect constraints on the charged Higgs mass in type II 2HDM [138]:

$$m_{H^+} > 240\ \mathrm{GeV} \quad \text{at } 95\%\ \mathrm{C.L.} \quad (41)$$

2HDM type III is a general model where both Higgs couple to up and down quarks. Of course, this means that 2HDM type III can induce flavor violation in the neutral current and thus it can be used to strongly constrain the new parameters in the model. We shall focus our interest on the two-Higgs doublet of type III, as the other two can be obtained from type III by taking some limits.

As in our previous sections for models of physics beyond the Standard Model, our strategy is to compute the Wilson coefficients contributing to the processes under consideration in the two-Higgs doublet model of type III, to check the consistency of our results with previous studies on these models and to use constraints coming from related D processes [139–142] in order to get an estimate of the maximum *CP* asymmetry that can be generated in the $D \to K\pi$ channel.

In the 2HDM of type III, the Yukawa Lagrangian can be written as [139,140]

$$\mathcal{L}_Y^{\mathrm{eff}} = \bar{Q}_{fL}^a[Y_{fi}^d \epsilon_{ab} H_d^{b\star} - \epsilon_{fi}^d H_u^a]d_{iR} - \bar{Q}_{fL}^a[Y_{fi}^u \epsilon_{ab} H_u^{b\star} + \epsilon_{fi}^u H_d^a]u_{iR} + \mathrm{H.c.}, \quad (42)$$

where ϵ_{ab} is the totally antisymmetric tensor, and ϵ_{ij}^q parametrizes the nonholomorphic corrections which couple up (down) quarks to the down (up) type Higgs doublet. After electroweak symmetry breaking, $\mathcal{L}_Y^{\mathrm{eff}}$ gives rise to the following charged Higgs-quark interaction Lagrangian:

$$\mathcal{L}_{H^\pm}^{\mathrm{eff}} = \bar{u}_f \Gamma_{u_f d_i}^{H^\pm LR\mathrm{eff}} P_R d_i + \bar{u}_f \Gamma_{u_f d_i}^{H^\pm RL\mathrm{eff}} P_L d_i, \quad (43)$$

with [140]

$$\Gamma_{u_f d_i}^{H^\pm LR\mathrm{eff}} = \sum_{j=1}^{3} \sin\beta V_{fj}\left(\frac{m_{d_i}}{v_d}\delta_{ji} - \epsilon_{ji}^d \tan\beta\right),$$
$$\Gamma_{u_f d_i}^{H^\pm RL\mathrm{eff}} = \sum_{j=1}^{3} \cos\beta \left(\frac{m_{u_f}}{v_u}\delta_{jf} - \epsilon_{jf}^{u\star}\tan\beta\right)V_{ji}. \quad (44)$$

Here v_u and v_d are the vacuum expectation values of the neutral component of the Higgs doublets, V is the CKM matrix, and $\tan\beta = v_u/v_d$. Using the Feynman rule given in Eq. (43), we derive the effective Hamiltonian resulting from the tree-level exchanging charged Higgs diagram that governs the process under consideration. We can express the effective Hamiltonian as

$$\mathcal{H}_{\mathrm{eff}} = \frac{G_F}{\sqrt{2}} V_{cs}^* V_{ud} \sum_{i=1}^{4} C_i^H(\mu) Q_i^H(\mu), \quad (45)$$

where C_i^H are the Wilson coefficients obtained by perturbative QCD running from $M_{H^\pm}$ scale to the scale μ relevant for hadronic decay, and Q_i^H are the relevant local operators at low energy scale $\mu \simeq m_c$. The operators can be written as

$$Q_1^H = (\bar{s}P_R c)(\bar{u}P_L d), \qquad Q_2^H = (\bar{s}P_L c)(\bar{u}P_R d),$$
$$Q_3^H = (\bar{s}P_L c)(\bar{u}P_L d), \qquad Q_4^H = (\bar{s}P_R c)(\bar{u}P_R d), \quad (46)$$

and for the Wilson coefficients C_i^H, at the electroweak scale, they can be expressed as

$$C_1^H = \frac{\sqrt{2}}{G_F V_{cs}^* V_{ud} m_H^2}\left(\sum_{j=1}^{3} \cos\beta V_{j1}\left(\frac{m_u}{v_u}\delta_{j1} - \epsilon_{j1}^{u\star}\tan\beta\right)\right)\left(\sum_{k=1}^{3}\cos\beta V_{k2}^{\star}\left(\frac{m_c}{v_u}\delta_{k2} - \epsilon_{k2}^{u}\tan\beta\right)\right),$$
$$C_2^H = \frac{\sqrt{2}}{G_F V_{cs}^* V_{ud} m_H^2}\left(\sum_{j=1}^{3} \sin\beta V_{1j}\left(\frac{m_d}{v_d}\delta_{j1} - \epsilon_{j1}^{d}\tan\beta\right)\right)\left(\sum_{k=1}^{3}\sin\beta V_{2k}^{\star}\left(\frac{m_s}{v_d}\delta_{k2} - \epsilon_{k2}^{d\star}\tan\beta\right)\right)$$
$$C_3^H = \frac{\sqrt{2}}{G_F V_{cs}^* V_{ud} m_H^2}\left(\sum_{j=1}^{3} \cos\beta V_{j1}\left(\frac{m_u}{v_u}\delta_{j1} - \epsilon_{j1}^{u\star}\tan\beta\right)\right)\left(\sum_{k=1}^{3}\sin\beta V_{2k}^{\star}\left(\frac{m_s}{v_d}\delta_{k2} - \epsilon_{k2}^{d\star}\tan\beta\right)\right),$$
$$C_4^H = \frac{\sqrt{2}}{G_F V_{cs}^* V_{ud} m_H^2}\left(\sum_{k=1}^{3} \cos\beta V_{k2}^{\star}\left(\frac{m_c}{v_u}\delta_{k2} - \epsilon_{k2}^{u}\tan\beta\right)\right)\left(\sum_{j=1}^{3}\sin\beta V_{1j}\left(\frac{m_d}{v_d}\delta_{j1} - \epsilon_{j1}^{d}\tan\beta\right)\right). \tag{47}$$

Having deriving the effective Hamiltonian which governs the process under consideration, we proceed to discuss the experimental constraints on the parameters ϵ_{ij}^q, where $q = d, u$, relevant to our decay mode. The flavor-changing elements ϵ_{ij}^d for $i \neq j$ are strongly constrained from flavor-changing neutral current (FCNC) processes in the down sector because of tree-level neutral Higgs exchange. Thus, we are left with only ϵ_{11}^d, ϵ_{22}^d. Concerning the elements ϵ_{ij}^u we see that only ϵ_{11}^u, ϵ_{22}^u can significantly affect the Wilson coefficients without any CKM suppression. Other ϵ_{ij}^u terms will be so small that the CKM suppression will be of orders λ or λ^2 or higher, and so we neglect them in our analysis. One of the important constraints on ϵ_{ij}^q, where $q = d, u$, can be obtained by applying the naturalness criterion of 't Hooft to the quark masses. According to the naturalness criterion of 't Hooft, the smallness of a quantity is only natural if a symmetry is gained in the limit in which this quantity is zero [140]. Thus, it is unnatural to have large accidental cancellations without a symmetry forcing these cancellations. Applying the naturalness criterion of 't Hooft to the quark masses in the 2HDM of type III, we find that [140]

$$|v_{u(d)}\epsilon_{ij}^{d(u)}| \leq |V_{ij}| \max[m_{d_i(u_i)}, m_{d_j(u_j)}], \tag{48}$$

which leads to

$$|\epsilon_{ij}^{d(u)}| \leq \frac{|V_{ij}| \max[m_{d_i(u_i)}, m_{d_j(u_j)}]}{|v_{u(d)}|}. \tag{49}$$

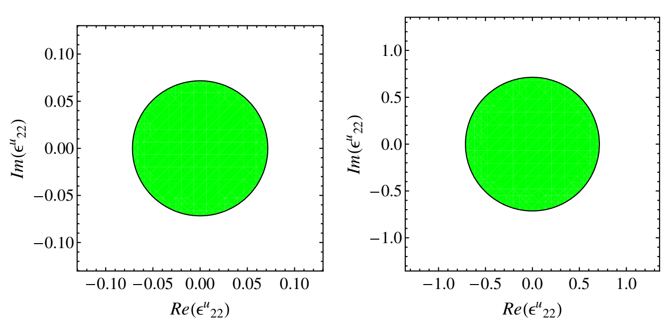

FIG. 3 (color online). Constraints on ϵ_{22}^u. Left: $\tan\beta = 10$. Right: $\tan\beta = 100$.

Clearly from the previous equation, ϵ_{11}^u, ϵ_{11}^d, ϵ_{22}^d will be severely constrained by their small masses while ϵ_{22}^u will be less constrained. Clearly from Eq. (49), the constraints imposed on ϵ_{22}^u are $\tan\beta$ dependent. We now apply the constraints imposed on the real and imaginary parts of ϵ_{22}^u corresponding to two different values of $\tan\beta$, namely, for two cases $\tan\beta = 10$ and $\tan\beta = 100$, using Eq. (49). In Fig. 3 we show the allowed regions for the two cases. Clearly the constraints are sensitive to the value of $\tan\beta$ where the constraints are weak for large values of $\tan\beta$. Since C_1^H and C_4^H are proportional to ϵ_{22}^u, they will be several orders of magnitude larger than C_2^H and C_3^H. In fact, this conclusion can be seen from Eq. (47) and thus in our analysis we drop C_2^H and C_3^H. Now, possible other constraints on ϵ_{22}^u can be obtained from $D - \bar{D}$ mixing and $K - \bar{K}$ mixing. For $K - \bar{K}$ mixing, the new contribution from the charged Higgs mediation corresponding to a top quark running in the loop will be much more dominant than the contribution in the case with the charm quark running in the loop. This is due to the dependency of the contribution on the ratio of the quark mass running in the loop to the charged Higgs mass. Thus, the expected constraints from $K - \bar{K}$ mixing might be relevant on ϵ_{32}^u and ϵ_{31}^u but not on ϵ_{22}^u. In fact, as mentioned in Ref. [140], the constraints on ϵ_{32}^u and ϵ_{31}^u are even weak and ϵ_{32}^u and ϵ_{31}^u can be sizable. By a similar argument we can neither use the process $b \to s\gamma$ nor the electric dipole moment (EDM) to constrain ϵ_{22}^u. Regarding $D - \bar{D}$ mixing, one expects a similar situation to that in $K - \bar{K}$ about the dominance of the top quark contribution. However, due to the CKM suppression factors, the top quark contribution will be smaller than the charm contribution.

1. $D - \bar{D}$ mixing constraints

In the following we discuss the possible constraints on ϵ_{22}^u from $D - \bar{D}$ mixing. We start by deriving the effective Hamiltonian that contributes to $D - \bar{D}$ mixing due to diagrams with charged Higgs mediation. In our calculations, we take into account only box diagrams that contribute to $D - \bar{D}$ mixing mediated by exchanging a strange quark and a charged Higgs. Other contributions from box diagrams mediated by down or bottom quarks and charged

Higgs are suppressed by the CKM factors. Since the SM contribution to $D-\bar{D}$ mixing is very small, we neglect its contribution and its interference with charged Higgs mediation contribution. Thus, the effective Hamiltonian for this case can be written as

$$\mathcal{H}_{H^\pm}^{|\Delta C|=2} = \frac{1}{m_{H^\pm}^2}\sum_{i=1}^{4} C_i(\mu)Q_i(\mu) + \tilde{C}_i(\mu)\tilde{Q}_i(\mu), \tag{50}$$

where C_i, $\tilde{C}_i$ are the Wilson coefficients obtained by perturbative QCD running from M_H scale to the scale μ relevant for hadronic decay and Q_i, $\tilde{Q}_i$ are the relevant local operators at the low energy scale

$$Q_1 = (\bar{u}\gamma^\mu P_L c)(\bar{u}\gamma_\mu P_L c), \qquad Q_2 = (\bar{u}P_L c)(\bar{u}P_L c),$$
$$Q_3 = (\bar{u}\gamma^\mu P_L c)(\bar{u}\gamma_\mu P_R c), \qquad Q_4 = (\bar{u}P_L c)(\bar{u}P_R c), \tag{51}$$

where we drop color indices and the operators $\tilde{Q}_i$ can be obtained from Q_i by changing the chirality $L \leftrightarrow R$. The Wilson coefficients C_i, are given by

$$C_1 = \frac{I_1(x_s)}{64\pi^2}\left(\sum_{j=1}^{3}\sin\beta V_{2j}^*\left(\frac{m_s}{v_d}\delta_{j2} - \epsilon_{j2}^d\tan\beta\right)\right)^2\left(\sum_{k=1}^{3}\sin\beta V_{1k}\left(\frac{m_s}{v_d}\delta_{k2} - \epsilon_{k2}^d\tan\beta\right)\right)^2,$$

$$C_2 = \frac{m_s^2 I_2(x_s)}{16\pi^2 m_{H^\pm}^2}\left(\sum_{j=1}^{3}\sin\beta V_{2j}^*\left(\frac{m_s}{v_d}\delta_{j2} - \epsilon_{j2}^d\tan\beta\right)\right)^2\left(\sum_{k=1}^{3}\cos\beta V_{k2}\left(\frac{m_u}{v_u}\delta_{k1} - \epsilon_{k1}^{u\star}\tan\beta\right)\right)^2,$$

$$C_3 = \frac{I_1(x_s)}{64\pi^2}\left(\sum_{j=1}^{3}\sin\beta V_{2j}^*\left(\frac{m_s}{v_d}\delta_{j2} - \epsilon_{j2}^d\tan\beta\right)\right)\left(\sum_{k=1}^{3}\sin\beta V_{1k}\left(\frac{m_s}{v_d}\delta_{k2} - \epsilon_{k2}^d\tan\beta\right)\right)\left(\sum_{l=1}^{3}\cos\beta V_{l2}\left(\frac{m_u}{v_u}\delta_{l1} - \epsilon_{l1}^{u\star}\tan\beta\right)\right)$$
$$\times\left(\sum_{n=1}^{3}\cos\beta V_{n2}^*\left(\frac{m_c}{v_u}\delta_{n2} - \epsilon_{n2}^{u\star}\tan\beta\right)\right),$$

$$C_4 = \frac{m_s^2 I_2(x_s)}{16\pi^2 m_{H^\pm}^2}\left(\sum_{j=1}^{3}\sin\beta V_{2j}^*\left(\frac{m_s}{v_d}\delta_{j2} - \epsilon_{j2}^d\tan\beta\right)\right)\left(\sum_{k=1}^{3}\sin\beta V_{1k}\left(\frac{m_s}{v_d}\delta_{k2} - \epsilon_{k2}^d\tan\beta\right)\right)\left(\sum_{l=1}^{3}\cos\beta V_{l2}\left(\frac{m_u}{v_u}\delta_{l1} - \epsilon_{l1}^{u\star}\tan\beta\right)\right)$$
$$\times\left(\sum_{n=1}^{3}\cos\beta V_{n2}^*\left(\frac{m_c}{v_u}\delta_{n2} - \epsilon_{n2}^{u\star}\tan\beta\right)\right), \tag{52}$$

where $x_s = m_s^2/m_{H^\pm}^2$, and the integrals are defined as follows:

$$I_1(x_s) = \frac{x_s+1}{(x_s-1)^2} + \frac{-2x_s\ln(x_s)}{(x_s-1)^3},$$
$$I_2(x_s) = \frac{-2}{(x_s-1)^2} + \frac{(x_s+1)\ln(x_s)}{(x_s-1)^3}. \tag{53}$$

The Wilson coefficients $\tilde{C}_i$ are given by

$$\tilde{C}_1 = \frac{I_1(x_s)}{64\pi^2}\left(\sum_{j=1}^{3}\cos\beta V_{j2}\left(\frac{m_u}{v_u}\delta_{j1} - \epsilon_{j1}^{u\star}\tan\beta\right)\right)^2$$
$$\times\left(\sum_{k=1}^{3}\cos\beta V_{k2}^*\left(\frac{m_c}{v_u}\delta_{k2} - \epsilon_{k2}^{u\star}\tan\beta\right)\right)^2,$$

$$\tilde{C}_2 = \frac{m_s^2 I_2(x_s)}{16\pi^2 m_{H^\pm}^2}\left(\sum_{j=1}^{3}\cos\beta V_{j2}^*\left(\frac{m_c}{v_u}\delta_{j2} - \epsilon_{j2}^{u\star}\tan\beta\right)\right)^2$$
$$\times\left(\sum_{k=1}^{3}\sin\beta V_{1k}\left(\frac{m_s}{v_d}\delta_{k2} - \epsilon_{k2}^d\tan\beta\right)\right)^2,$$

$$\tilde{C}_3 = C_3, \qquad \tilde{C}_4 = C_4. \tag{54}$$

Our operators Q_1, Q_2, and Q_4 given in Eq. (51) are equivalent to their corresponding operators given in Refs. [141,142], while the operators $\tilde{Q}_1$ and $\tilde{Q}_2$ are equivalent to Q_6 and Q_7 given in the same references, respectively. Moreover, Q_3, given in Eq. (51), can be related to Q_5 in Refs. [141,142] by Fierz identity. For the rest of the operators, $\tilde{Q}_3$ and $\tilde{Q}_4$, they are equivalent to Q_5 and Q_4 in Refs. [141,142] since their matrix elements are equal. Thus, the set of the Wilson coefficients C_i and $\tilde{C}_i$, for i running from 1 to 4 that we derived above, are subjected to the constraints given in Refs. [141,142], which in our case read

$$|C_1| \le 5.7\times 10^{-7}\left[\frac{m_{H^\pm}}{1\ \text{TeV}}\right]^2$$
$$|C_2| \le 1.6\times 10^{-7}\left[\frac{m_{H^\pm}}{1\ \text{TeV}}\right]^2$$
$$|C_3| \le 3.2\times 10^{-7}\left[\frac{m_{H^\pm}}{1\ \text{TeV}}\right]^2$$
$$|C_4| \le 5.6\times 10^{-8}\left[\frac{m_{H^\pm}}{1\ \text{TeV}}\right]^2. \tag{55}$$

The constraints on $\tilde{C}_1 - \tilde{C}_4$ are similar to those on $C_1 - C_4$. As can be seen from Eq. (55), the constraints on the Wilson coefficients will be strong for small charged Higgs masses. With the explicit expressions for the Wilson coefficients given above, we can proceed now to derive the constraints on ϵ_{22}^u using the upper bound on $\tilde{C}_2$, for instance. Keeping terms corresponding to first order in λ, where λ is the CKM parameter, we find that, for $m_{H^\pm} = 300$ GeV and $\tan\beta = 55$,

$$\tilde{C}_2 \times 10^{12} \simeq 3(-53.6\epsilon^d_{12} - 12.7\epsilon^d_{22} + 0.007)^2 \times (-12.4\epsilon^{u*}_{12} - 53.4\epsilon^{u*}_{22} + 0.007)^2. \quad (56)$$

While for $m_{H^\pm} = 300$ GeV and $\tan\beta = 500$, we find

$$\tilde{C}_2 \times 10^{14} \simeq 3.6(-487.1\epsilon^d_{12} - 115.0\epsilon^d_{22} + 0.06)^2 \times (-112.5\epsilon^{u*}_{12} - 486.7\epsilon^{u*}_{22} + 0.007)^2. \quad (57)$$

In both Eqs. (57) and (56) we can drop terms proportional to ϵ^{u*}_{12} to a good approximation, as they have small coefficients in comparison to ϵ^u_{22}, and also since $\epsilon^{u,d}_{ij}$ with $i \neq j$ are always smaller than the diagonal elements $\epsilon^{u,d}_{ii}$. On the other hand, we know that ϵ^d_{12} cannot be large to not allow flavor-changing neutral currents, and so we can also drop terms proportional to ϵ^d_{12} in Eqs. (57) and (56) to a good approximation. Thus, we are left with ϵ^d_{22} and ϵ^u_{22} in both Eqs. (57) and (56). Comparing their coefficients shows that ϵ^u_{22} has a large coefficient and thus we can drop ϵ^d_{22} terms. An alternative way is to assume that ϵ^u_{22} terms are the dominant ones, in comparison to the other $\epsilon^{u,d}_{ij}$ terms, and proceed to set upper bounds on ϵ^u_{22}. In fact, even if we consider other Wilson coefficients rather than $\tilde{C}_2$, this conclusion will not be altered. Under the assumption $\epsilon^d_{12} = \epsilon^d_{22} = \epsilon^u_{12} = 0$ and using the upper bound corresponding to $m_{H^\pm} = 300$ GeV on $\tilde{C}_2$, given by Eq. (55), one obtains

$$|\tilde{C}_2| \leq 1.4 \times 10^{-8}. \quad (58)$$

Clearly from Eqs. (57)–(59) the bounds that can be obtained on ϵ^u_{22} will be very loose, and thus $D - \bar{D}$ mixing cannot lead to strong constraints on ϵ^u_{22}.

2. $D_q \to \tau\nu$ constraints

The decay modes $D_q \to \tau\nu$, where $q = d$ or $q = s$, can be generated in the SM at tree level via W boson mediation. Within the 2HDM of type III under consideration, the charged Higgs can mediate these decay modes at tree level also and hence, the total branching ratios, following a similar notation in Ref. [140], can be expressed as

$$\mathcal{B}(D^+_q \to \tau^+\nu) = \frac{G_F^2|V_{cq}|^2}{8\pi} m_\tau^2 f^2_{D_q} m_{D_q}\left(1 - \frac{m^2_\tau}{m^2_{D_q}}\right)^2 \tau_{D_q} \times \left|1 + \frac{m^2_{D_q}}{(m_c + m_q)m_\tau}\frac{(C^{cq*}_R - C^{cq*}_L)}{C^{cq*}_{\rm SM}}\right|^2, \quad (59)$$

where we have used [143]

$$\langle 0|\bar{q}\gamma^5 c|D_q\rangle = \frac{f_{D_q} m^2_{D_q}}{(m_c + m_q)}, \quad (60)$$

where the SM Wilson coefficient is given by $C^{cq}_{\rm SM} = 4G_F V_{cq}/\sqrt{2}$ and the Wilson coefficients C^{cq}_L and C^{cq}_R at the matching scale are given by

$$C^{cq}_{R(L)} = \frac{-1}{M^2_{H^\pm}}\Gamma^{LR(RL),H^\pm}_{cq}\frac{m_\tau}{v}\tan\beta, \quad (61)$$

with the vacuum expectation value $v \approx 174$ GeV and $\Gamma^{LR(RL),H^\pm}_{cq}$ can be read from Eq. (44). Setting the charged Higgs contribution to zero and $f_{D_s} = 248 \pm 2.5$ MeV [144], we find that $\mathcal{B}^{\rm SM}(D^+_d \to \tau^+\nu) \simeq 9.5 \times 10^{-4}$ and $\mathcal{B}^{\rm SM}(D^+_s \to \tau^+\nu) = (5.11 \pm 0.11) \times 10^{-2}$, which are in close agreement with the results in Refs. [145–147]. The experimental values of these branching ratios are given by $\mathcal{B}(D^+_d \to \tau^+\nu) < 2.1 \times 10^{-3}$ [148], while $\mathcal{B}(D^+_s \to \tau^+\nu) = (5.38 \pm 0.32) \times 10^{-2}$ [149]. Keeping the terms that are proportional to the dominant CKM elements, we find for $q = d$

$$\Gamma^{H^\pm RL{\rm eff}}_{cd} = \cos\beta V_{11}(-\epsilon^{u*}_{12}\tan\beta)$$
$$\Gamma^{H^\pm LR{\rm eff}}_{cd} = \sin\beta V_{11}\left(\frac{m_d}{v_d} - \epsilon^d_{11}\tan\beta\right), \quad (62)$$

while for $q = s$ we find

$$\Gamma^{H^\pm RL{\rm eff}}_{cs} = \cos\beta V_{22}\left(\frac{m_c}{v_u} - \epsilon^{u*}_{22}\tan\beta\right)$$
$$\Gamma^{H^\pm LR{\rm eff}}_{cs} = \sin\beta V_{22}\left(\frac{m_s}{v_d} - \epsilon^d_{22}\tan\beta\right). \quad (63)$$

Clearly from the last two equations, we need to consider the decay mode $D^+_s \to \tau^+\nu$ to constrain ϵ^u_{22}. For $\tan\beta = 10$ we find that

$$\Gamma^{H^\pm RL{\rm eff}}_{cs} \times 10^{-3} \simeq 0.71 - 968.6\epsilon^u_{22}$$
$$\Gamma^{H^\pm LR{\rm eff}}_{cs} \times 10^{-3} \simeq 5.3 - 9686.0\epsilon^d_{22}. \quad (64)$$

Clearly the coefficient of ϵ^d_{22} is 1 order of magnitude larger than ϵ^u_{22} and for larger $\tan\beta$ one expects it to be larger. However, ϵ^d_{22} is severely constrained by the naturalness criterion and we expect the term proportional to ϵ^u_{22} to be larger; in our analysis we can drop the ϵ^d_{22} term and proceed to obtain the required constraints. We show in Figs. 4 and 5 the allowed regions for the real and imaginary parts of ϵ^u_{22} corresponding to two different values of the charged Higgs mass, namely, $m_{H^\pm} = 150$ GeV and $m_{H^\pm} = 180$ GeV, and for different values of $\tan\beta$. Our objective here is to show the dependency of the constraints on $m_{H^\pm}$ and $\tan\beta$. We see from the figures that for $\tan\beta = 140$, the constraints become loose with increasing $m_{H^\pm}$. It is important to emphasize that such a large value for $\tan\beta$ is compatible with supersymmetry-inspired two-Higgs doublet models, as shown in Ref. [150]. This is expected, as Wilson coefficients of the charged Higgs are inversely proportional to the square of $m_{H^\pm}$ and thus their contributions to $\mathcal{B}(D^+_s \to \tau^+\nu)$ become small for large $m_{H^\pm}$, which in turn means that the constraints obtained are loose. Another remark from Fig. 4 is that, for a fixed Higgs mass, the constraints become strong with increasing $\tan\beta$, which is also expected from Eq. (61). This in contrast to the constraints derived by applying the naturalness criterion, where we showed that the constraints become loose with increasing $\tan\beta$.

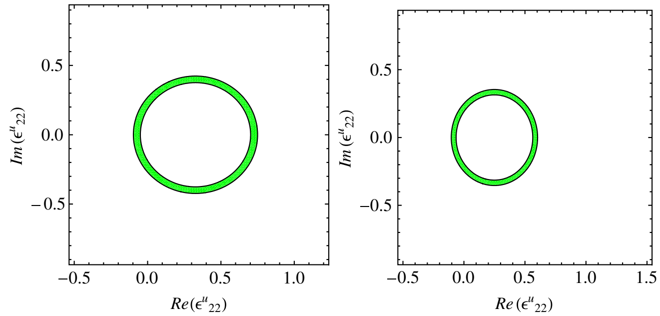

FIG. 4 (color online). Allowed 2σ regions for the real and imaginary parts of ϵ^u_{22} from $\mathcal{B}(D_s^+ \to \tau^+ \nu)$. Left: $\tan\beta = 117$. Right: $\tan\beta = 140$. In both cases we take $m_{H^\pm} = 150$ GeV.

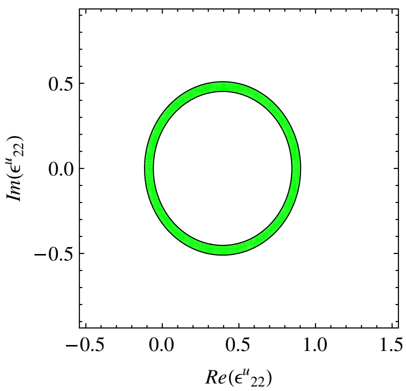

FIG. 5 (color online). Allowed 2σ regions for the real and imaginary parts of ϵ^u_{22} from $\mathcal{B}(D_s^+ \to \tau^+ \nu)$ corresponding to $\tan\beta = 140$ and $m_{H^\pm} = 180$ GeV.

3. *CP violation in charged Higgs*

The total amplitude including the SM and charged Higgs contribution can be written as

$$\mathcal{A} = \left(C_1^{\rm SM} + \frac{1}{N}C_2^{\rm SM} + \chi^{\pi^+}(C_1^H - C_4^H)\right)X^{\pi^+}_{D^0K^-} - \left(C_2^{\rm SM} + \frac{1}{N}C_1^{\rm SM} + \frac{1}{2N}(C_1^H - \chi^{D^0}C_4^H)\right)X^{D^0}_{K^-\pi^+}, \quad (65)$$

with $X^{P_1}_{P_2P_3} = i f_{P_1}\Delta^2_{P_2P_3}F_0^{P_2P_3}(m^2_{P_1})$, $\Delta^2_{P_2P_3} = m^2_{P_2} - m^2_{P_3}$, and χ^{π^+} and χ^{D^0} were previously defined as

$$\chi^{\pi^+} = \frac{m_\pi^2}{(m_c - m_s)(m_u + m_d)}$$
$$\chi^{D^0} = \frac{m_D^2}{(m_c + m_u)(m_s - m_d)}. \quad (66)$$

The form of the amplitude, $\mathcal{A}$, shows how a charged Higgs contribution can affect only the short distance physics (Wilson coefficients) without any new effect on the long range physics (hadronic parameters). Thus, the strong phase will not be affected by including charged Higgs contributions, while the weak phase will be affected. We can rewrite Eq. (65) in terms of the amplitudes T and E introduced before in the case of the SM as follows:

$$\mathcal{A} = V^*_{cs}V_{ud}(T^{\rm SM+H} + E^{\rm SM+H}), \quad (67)$$

where

$$T^{\rm SM+H} = 3.14\times 10^{-6} \simeq \frac{G_F}{\sqrt{2}}a_1^{\rm SM+H}f_\pi(m_D^2 - m_K^2)F_0^{DK}(m_\pi^2)$$
$$E^{\rm SM+H} = 1.53\times 10^{-6}e^{122^\circ i} \simeq \frac{G_F}{\sqrt{2}}a_2^{\rm SM+H}f_D(m_K^2 - m_\pi^2)F_0^{K\pi}(m_D^2), \quad (68)$$

where

$$a_1^{\rm SM+H} = \left(C_1^{\rm SM} + \frac{1}{N}C_2^{\rm SM} + \chi^{\pi^+}(C_1^H - C_4^H)\right) = (a_1 + \Delta a_1 + \chi^{\pi^+}(C_1^H - C_4^H)) \quad (69)$$

$$a_2^{\rm SM+H} = -\left(a_2 + \Delta a_2 + \frac{1}{2N}(C_1^H - \chi^{D^0}C_4^H)\right). \quad (70)$$

The *CP* asymmetry can be obtained using the relation

$$A_{CP} = \frac{|\mathcal{A}|^2 - |\bar{\mathcal{A}}|^2}{|\mathcal{A}|^2 + |\bar{\mathcal{A}}|^2} = \frac{2|T^{\rm SM+H}||E^{\rm SM+H}|\sin(\phi_1 - \phi_2)\sin(-\alpha_E)}{|T^{\rm SM+H} + E^{\rm SM+H}|^2}, \quad (71)$$

with $\phi_i = {\rm Arg}[a_i^{\rm SM+H}]$ and $\alpha_E = {\rm Arg}(\chi_E)$. As an example let us take ${\rm Re}(\epsilon^u_{22}) = 0.04$, ${\rm Im}(\epsilon^u_{22}) = 0.03$, which is an allowed point for $\tan\beta = 10$. In this case we find that for a value of $m_{H^\pm} = 500$ GeV we find that $A_{CP} \simeq -3.7\times 10^{-5}$, while for $m_H = 300$ GeV we find that $A_{CP} \simeq -1\times 10^{-4}$. Let us take another example where ${\rm Re}(\epsilon^u_{22}) = -0.5$, ${\rm Im}(\epsilon^u_{22}) = -0.3$, which is an allowed point for $\tan\beta = 100$. Repeating the same steps as above, we find that for $m_{H^\pm} = 250$ GeV the predicted $A_{CP} \simeq 1.5\times 10^{-2}$. Clearly in charged Higgs models the predicted *CP* asymmetry is very sensitive to the value of $\tan\beta$ and to the value of the Higgs mass.

V. CONCLUSION

In this paper, we have studied the Cabibbo favored nonleptonic D^0 decays into $K^-\pi^+$. We have shown that the Standard Model prediction for the corresponding *CP* asymmetry is strongly suppressed and out of experimental range, even when taking into account the large strong phases coming from the final state interactions. Then we explored new physics models, taking into account three possible extensions, namely, extra family, extra gauge bosons within left-right grand unification models, and extra Higgs fields. The fourth family model strongly improved the SM prediction of the *CP* asymmetry but still the predicted *CP* asymmetry is far out of the reach of LHCb or a SuperB factory such as SuperKEKB. The most promising models are nonmanifest left-right extensions of the SM where the LR mixing between the gauge bosons permits us to get a strong enhancement in the *CP* asymmetry. In such a model, it is possible to get *CP* asymmetry of order 10%, which is within the range of LHCb and the next generation of charm or B factories. The nonobservation of such a huge CP asymmetry will strongly constrain the parameters of this model. In multi-Higgs extensions of the SM, the 2HDM type III is the most attractive, as it permits us to solve the puzzle coming from $B \to \tau\nu$ and at the same time give a large contribution to this CP asymmetry depending on the charged Higgs masses and couplings. A maximal value of 1.5% can be reached with a Higgs mass of 300 GeV and large $\tan\beta$.

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APPENDIX: OPERATORS AND OTHER DEFINITIONS

We start by defining $X^{P_1}_{P_2P_3}$, where P_i denotes a pseudoscalar meson, as follows:

$$X^{P_1}_{P_2P_3} = if_{P_1}\Delta^2_{P_2P_3}F_0^{P_2P_3}(m^2_{P_1}), \tag{A1}$$

where $\Delta^2_{P_2P_3} = m^2_{P_2} - m^2_{P_3}$. In terms of $X^{P_1}_{P_2P_3}$ we find that

$$\begin{aligned}
&\langle\pi^+|\bar{u}\gamma_\mu\gamma_5 d|0\rangle\langle K^-|\bar{s}\gamma_\mu c|D^0\rangle = -X^{\pi^+}_{D^0K^-} \qquad \langle K^-\pi^+|\bar{s}\gamma_\mu d|0\rangle\langle 0|\bar{u}\gamma_\mu\gamma_5 c|D^0\rangle = X^{D^0}_{K^-\pi^+}\\
&\langle\pi^+|\bar{u}\gamma_5 d|0\rangle\langle K^-|\bar{s}c|D^0\rangle = -\frac{m^2_\pi}{(m_c-m_s)(m_u+m_d)}X^{\pi^+}_{D^0K^-} \equiv -\chi^{\pi^+}X^{\pi^+}_{D^0K^-}\\
&\langle K^-\pi^+|\bar{s}d|0\rangle\langle 0|\bar{u}\gamma_5 c|D^0\rangle = -\frac{m^2_D}{(m_c+m_u)(m_s-m_d)}X^{D^0}_{K^-\pi^+} \equiv -\chi^{D^0}X^{D^0}_{K^-\pi^+}.
\end{aligned} \tag{A2}$$

Using Eq. (A2) we get

$$\begin{aligned}
&\langle K^-\pi^+|\mathcal{O}_1|D^0\rangle = \langle K^-\pi^+|\bar{s}\gamma_\mu c_L\bar{u}\gamma_\mu d_L|D^0\rangle = \langle\pi^+|\bar{u}\gamma_\mu d_L|0\rangle\langle K^-|\bar{s}\gamma_\mu c_L|D^0\rangle\\
&\quad + \frac{1}{N}\langle K^-\pi^+|\bar{s}\gamma_\mu d_L|0\rangle\langle 0|\bar{u}\gamma_\mu c_L|D^0\rangle = X^{\pi^+}_{D^0K^-} - \frac{1}{N}X^{D^0}_{K^-\pi^+}\\
&\langle K^-\pi^+|\mathcal{O}_2|D^0\rangle = \langle K^-\pi^+|\bar{u}\gamma_\mu c_L\bar{s}\gamma_\mu d_L|D^0\rangle = \langle K^-\pi^+|\bar{s}\gamma_\mu d_L|0\rangle\langle 0|\bar{u}\gamma_\mu c_L|D^0\rangle\\
&\quad + \frac{1}{N}\langle\pi^+|\bar{u}\gamma_\mu d_L|0\rangle\langle K^-|\bar{s}\gamma_\mu c_L|D^0\rangle = -X^{D^0}_{K^-\pi^+} + \frac{1}{N}X^{\pi^+}_{D^0K^-}\\
&\langle K^-\pi^+|\bar{s}\gamma_\mu c_R\bar{u}\gamma_\mu d_R|D^0\rangle = \langle\pi^+|\bar{u}\gamma_\mu d_R|0\rangle\langle K^-|\bar{s}\gamma_\mu c_R|D^0\rangle + \frac{1}{N}\langle K^-\pi^+|\bar{s}\gamma_\mu d_R|0\rangle\langle 0|\bar{u}\gamma_\mu c_R|D^0\rangle = -\langle K^-\pi^+|\mathcal{O}_1|D^0\rangle\\
&\langle K^-\pi^+|\bar{u}\gamma_\mu c_R\bar{s}\gamma_\mu d_R|D^0\rangle = \langle K^-\pi^+|\bar{s}\gamma_\mu d_R|0\rangle\langle 0|\bar{u}\gamma_\mu c_R|D^0\rangle + \frac{1}{N}\langle\pi^+|\bar{u}\gamma_\mu d_R|0\rangle\langle K^-|\bar{s}\gamma_\mu c_R|D^0\rangle = -\langle K^-\pi^+|\mathcal{O}_2|D^0\rangle\\
&\langle K^-\pi^+|\bar{s}\gamma_\mu c_L\bar{u}\gamma_\mu d_R|D^0\rangle = \langle\pi^+|\bar{u}\gamma_\mu d_R|0\rangle\langle K^-|\bar{s}\gamma_\mu c_L|D^0\rangle - \frac{2}{N}\langle K^-\pi^+|\bar{s}d_{S+P}|0\rangle\langle 0|\bar{u}c_{S-P}|D^0\rangle = -X^{\pi^+}_{D^0K^-} - \frac{2}{N}\chi^{D^0}X^{D^0}_{K^-\pi^+}\\
&\langle K^-\pi^+|\bar{u}\gamma_\mu c_L\bar{s}\gamma_\mu d_R|D^0\rangle = \langle K^-\pi^+|\bar{s}\gamma_\mu d_R|0\rangle\langle 0|\bar{u}\gamma_\mu c_L|D^0\rangle - \frac{2}{N}\langle\pi^+|\bar{u}d_{S+P}|0\rangle\langle K^-|\bar{s}c_{S-P}|D^0\rangle = -X^{D^0}_{K^-\pi^+} + \frac{2}{N}\chi^{\pi^+}X^{\pi^+}_{D^0K^-}\\
&\langle K^-\pi^+|\bar{s}\gamma_\mu c_R\bar{u}\gamma_\mu d_L|D^0\rangle = \langle\pi^+|\bar{u}\gamma_\mu d_L|0\rangle\langle K^-|\bar{s}\gamma_\mu c_R|D^0\rangle - \frac{2}{N}\langle K^-\pi^+|\bar{s}d_{S-P}|0\rangle\langle 0|\bar{u}c_{S+P}|D^0\rangle = X^{\pi^+}_{D^0K^-} + \frac{2}{N}\chi^{D^0}X^{D^0}_{K^-\pi^+}\\
&\langle K^-\pi^+|\bar{u}\gamma_\mu c_R\bar{s}\gamma_\mu d_L|D^0\rangle = \langle K^-\pi^+|\bar{s}\gamma_\mu d_L|0\rangle\langle 0|\bar{u}\gamma_\mu c_R|D^0\rangle - \frac{2}{N}\langle\pi^+|\bar{u}d_{S-P}|0\rangle\langle K^-|\bar{s}c_{S+P}|D^0\rangle = X^{D^0}_{K^-\pi^+} - \frac{2}{N}\chi^{\pi^+}X^{\pi^+}_{D^0K^-},
\end{aligned} \tag{A3}$$

and for the scalar ones

$$
\begin{aligned}
&\langle K^-\pi^+|\bar{s}c_L\bar{u}d_L|D^0\rangle = \chi^{\pi^+}X^{\pi^+}_{D^0K^-} - \frac{1}{2N}\chi^{D^0}X^{D^0}_{K^-\pi^+} \qquad &&\langle K^-\pi^+|\bar{u}c_L\bar{s}d_L|D^0\rangle = \chi^{D^0}X^{D^0}_{K^-\pi^+} - \frac{1}{2N}\chi^{\pi^+}X^{\pi^+}_{D^0K^-}\\
&\langle K^-\pi^+|\bar{s}c_R\bar{u}d_R|D^0\rangle = -\chi^{\pi^+}X^{\pi^+}_{D^0K^-} + \frac{1}{2N}\chi^{D^0}X^{D^0}_{K^-\pi^+} \qquad &&\langle K^-\pi^+|\bar{u}c_R\bar{s}d_R|D^0\rangle = -\chi^{D^0}X^{D^0}_{K^-\pi^+} + \frac{1}{2N}\chi^{\pi^+}X^{\pi^+}_{D^0K^-}\\
&\langle K^-\pi^+|\bar{s}c_L\bar{u}d_R|D^0\rangle = -\chi^{\pi^+}X^{\pi^+}_{D^0K^-} + \frac{1}{2N}X^{D^0}_{K^-\pi^+} \qquad &&\langle K^-\pi^+|\bar{u}c_L\bar{s}d_R|D^0\rangle = \chi^{D^0}X^{D^0}_{K^-\pi^+} + \frac{1}{2N}X^{\pi^+}_{D^0K^-}\\
&\langle K^-\pi^+|\bar{s}c_R\bar{u}d_L|D^0\rangle = \chi^{\pi^+}X^{\pi^+}_{D^0K^-} - \frac{1}{2N}X^{D^0}_{K^-\pi^+} \qquad &&\langle K^-\pi^+|\bar{u}c_R\bar{s}d_L|D^0\rangle = -\chi^{D^0}X^{D^0}_{K^-\pi^+} - \frac{1}{2N}X^{\pi^+}_{D^0K^-},
\end{aligned} \tag{A4}
$$

where Fierz's ordering has been used,

$$
\begin{aligned}
&(\bar{\psi}_1\Psi_2)_L(\bar{\psi}_3\Psi_4)_L = (\bar{\psi}_1\Psi_4)_L(\bar{\psi}_3\Psi_2)_L, \qquad (\bar{\psi}_1\Psi_2)_L(\bar{\psi}_3\Psi_4)_R = -2(\bar{\psi}_1\Psi_4)_{S+P}(\bar{\psi}_3\Psi_2)_{S-P}\\
&4\bar{\psi}_1\psi_{2,S\pm P}\bar{\psi}_3\psi_{4,S\pm P} = -2\bar{\psi}_1\psi_{4,S\pm P}\bar{\psi}_3\psi_{2,S\pm P} - \frac{1}{2}\bar{\psi}_1(1\pm\gamma_5)\sigma^{\mu\nu}\psi_4\bar{\psi}_3(1\pm\gamma_5)\sigma_{\mu\nu}\psi_2\\
&2(T_a)_{\alpha\beta}(T_a)_{\gamma\delta} = \delta_{\alpha\delta}\delta_{\beta\gamma} - \frac{1}{N}\delta_{\alpha\beta}\delta_{\gamma\delta}, \qquad (T_a)_{\alpha\beta}(T_a)_{\gamma\delta} = \frac{N_C^2-1}{2N_C}\delta_{\alpha\delta}\delta_{\beta\gamma} - \frac{1}{N_C}(T^a)_{\alpha\delta}(T^a)_{\beta\gamma}.
\end{aligned} \tag{A5}
$$

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