Observation of CP violation in $D^0 \rightarrow K^-\pi^+$ as a smoking gun for New Physics

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Abstract

In this paper, we study the Cabibbo favored non-leptonic $D^0$ decays into $K^-\pi^+$ decays. First we show that, within the Standard Model, the corresponding 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 few orders of magnitude however the resulting CP asymmetry is still far from experimental range. The most sensitive New Physics Models to this CP asymmetry comes from no-manifest Left-Right models where a CP asymmetry up to 10% can be reached and general two Higgs models extension of 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, mixings angles and CP violating phases appearing in the Cabibbo-Kobayashi-Maskawa (CKM) quarks 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-quarks sector and even tensions with experimental measurements remain to be cleared (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. One of these processes is the rare decays and other ‘null’ tests which correspond to an observable strictly equal to zero within SM. So any deviation from zero of these ‘null’ tests observables is a clear signal of Physics beyond SM. This is the case of Cabibbo-Favored (CF) and Double Cabibbo Suppressed (DCS) non-leptonic 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 (SCD) modes)\[17–32\], it is not clear that the SM\[33–41\] can describe correctly the CP violation in the up quarks sector. It is even more difficult as large distance contributions are important and difficult to be evaluated\[42–46\]. From the point of view of New Physics (NP), CP violation in CF and DCS modes is an excellent opportunity given that it is very suppressed in the SM and it is not easy to find a NP model able to produce a reasonable CP violation signal. Thus measuring CP violation in these channels is a very clear signal of New Physics.

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

\[ x \equiv \frac{\Delta m_d}{\Gamma_D} = 0.55^{+0.12}_{-0.13}, \quad y \equiv \frac{\Delta \Gamma_D}{2\Gamma_D} = 0.83(13) \]
\[ , \quad q \equiv \frac{\Delta m_d}{\Gamma_D} = 0.91^{+0.18}_{-0.16}, \quad \phi \equiv \arg \left( \frac{q}{p} \right) = - \left( 10.2^{+9.4}_{-8.9} \right)^\circ \]
| Mode                  | BR[%]     | $A_{\text{CP}}$ [%] | Mode                  | BR[%]     | $A_{\text{CP}}$ [%] |
|----------------------|-----------|---------------------|----------------------|-----------|---------------------|
| $D^0 \to K^-\pi^+$ CF | 3.95(5)   | -                   | $D^0 \to K^0\pi^0$ 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 | $1.48(7) \cdot 10^{-4}$ | -                   | $D^0 \to K^0\pi^0$ 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 | $1.72(19) \cdot 10^{-2}$ | -                   |
| $D^+ \to K^+\eta$ DCS | $1.08(17) \cdot 10^{-2}$ | -                   | $D^+ \to K^+\eta'$ DCS | $1.76(22) \cdot 10^{-2}$ | -                   |
| $D^+_s \to K^+K^0$ DCS | -         | -                   |
| $D^0 \to \pi^-\pi^+$ | $0.143(3)$ | $0.22(24)(11)$      | $A_{\text{CP}}(K^+K^-) - A_{\text{CP}}(\pi^+\pi^-)$ | -         | $-0.65(18)$         |
| $D^0 \to K^-K^+$ | $0.398(7)$ | $-0.24(22)(9)$      | $A_{\text{CP}}(K^+\bar{K}^0) - A_{\text{CP}}(\bar{K}^0\bar{K}^0)$ | -         | -                   |
| $D^+ \to K^0\pi^+$ | $1.47(7)$ | $-0.71(19)(20)$     | $D^+ \to \pi^+\pi^-\pi^\pm$ | $0.327(22)$ | $1.7(42)$           |
| $D^\pm \to K^0\pi^\pm\pi^\mp$ | $9.51(34)$ | $-0.5(4)(9)$       | $D^\mp \to K^0\pi^\mp\pi^0$ | $6.90(32)$ | $0.3(9)(3)$        |
| $D^\pm \to K^\mp K^-\pi^\pm$ | $0.98(4)$ | $0.39(61)$         |

TABLE I. Direct CP in $D$ non-leptonic decays, from Heavy Flavor Averaging Group HAFG[1, 51]

where $x \neq 0$ or/and $y \neq 0$ 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 theory (HQE) (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 GIM 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 factors: CKM matrix elements ($|V_{ub}V_{cb}^*/V_{us}V_{cs}^*| \sim 10^{-6}$) and the fact that $b$ quark mass is small compared to top mass. CP violation in the $b$-quark sector is due to the large top quark mass, while in the kaon 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 project [50]. In table [1] the experimentally measured Branching ratios and CP asymmetries are given for different non-leptonic $D$ decays. In this paper, we study in details the CP asymmetry for the CF $D^0 \to K^-\pi^+$ decay.
In sect. II, we give the general description of the Effective Hamiltonian describing this decay within 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 sect. 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 sect. IV, New Physics models are introduced and their contributions to CP asymmetry are evaluated. Finally, we conclude in sect. V.

II. GENERAL DESCRIPTION OF CF NON LEPTONIC $D^0$ DECAYS INTO $K^-$ AND $\pi^+$

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

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

(3)

with $i =$S, V and T for respectively scalar (S), vectorial (V) and tensorial (T) operators. 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.

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

(4)

$$= \frac{G_F}{\sqrt{2}} V_{cs}^* V_{ud} \left( c_1 \mathcal{O}_1 + c_2 \mathcal{O}_2 \right) + \text{h.c.}$$

(5)

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

$$A_{D^0 \rightarrow K^- \pi^+} = -\frac{G_F}{\sqrt{2}} V_{cs}^* V_{ud} \left[ a_1 X_{D^0 K^-}^{\pi^+} + a_2 X_{K^- \pi^+}^{D^0} \right],$$

(6)

$$\text{BR} = \frac{\tau_D p_K}{8\pi m_D^2} |A|^2$$

(7)

where BR is the Branching ratio of the process. $\tau_D$ is the D lifetime, $p_K$ is the Kaon momentum and $m_D$ is the D meson mass. The $X_{D^0 K^-}^{\pi^+}$ and $X_{K^- \pi^+}^{D^0}$ can be expressed in the following way:

$$X_{P_2 P_3} = i f_{P_1} \Delta P_2 P_3 F_0 (m_{P_1}^2), \ \Delta_{P_2 P_3}^2 = m_{P_2}^2 - m_{P_3}^2$$

(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 Final State Interaction (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 you have many FSI contributions: resonances, other intermediate states, rescattering, and so on. Resonances are specially important in this region given that they are abundant. They can be included and seems 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 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), exchange quark diagrams (E), color-suppressed quarks diagrams (C). Their results can be summarized in the following way, for the process under consideration \([56]\):

\[
A_{D^0 \rightarrow K^- \pi^+} \equiv V_{cs}^* V_{ud} (T + E)
\]  

(9)

with

\[
T = (3.14 \pm 0.06) \cdot 10^{-6} \text{GeV} \\
E = 1.53_{-0.08}^{+0.07} \cdot 10^{-6} \cdot e^{(122\pm2)\circ i} \text{GeV}
\]

(10)

where in NFA they can be approximately written as

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

(11)

\[
E \simeq -\frac{G_F}{\sqrt{2}} a_2 f_D (m_K^2 - m_\pi^2) F_0^{K\pi}(m_D^2)
\]

(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 \( D^0 \rightarrow K^- \pi^+ \) WITHIN SM

In the case of CF (and DCF) processes the corrections are very small (see diagrams in fig.1 and fig.2) and are generated through box and di-penguin diagrams \([57, 59]\). In this section, we shall evaluate these contributions.
The box contribution is given as \[59, 60\]

\[
\Delta H = \frac{G_F^2 m_W^2}{2\pi^2} V^*_{cd} V_{us} V^*_{ud} f(x_U, x_D) \bar{u} \gamma_\mu c_L \bar{s} \gamma^\mu d_L
\]

\[
= \frac{G_F^2 m_W^2}{2\pi^2} \lambda^D_{cd} \lambda^U_{sd} f(x_U, x_D) O_2
\]

\[
= \frac{G_F^2 m_W^2}{2\pi^2} b_x O_2
\]

where

\[
b_x \equiv \lambda^D_{cu} \lambda^U_{sd} f(x_U, x_D)
\]

\[
= V^*_{cd} V_{ud} (V^*_{us} V_{ud} f_{us} + V^*_{cs} V_{cd} f_{cd} + V^*_{ts} V_{td} f_{td})
\]

\[
+ V^*_{cs} V_{us} (V^*_{us} V_{ud} f_{us} + V^*_{cs} V_{cd} f_{cd} + V^*_{ts} V_{td} f_{td}) + V^*_{cb} V_{ub} (V^*_{us} V_{ud} f_{ub} + V^*_{cs} V_{cd} f_{cb} + V^*_{ts} V_{td} f_{tb})
\]

\[
= V^*_{cs} V_{us} [V^*_{cs} V_{cd} (f_{cs} - f_{cd} - f_{us} + f_{ud}) + V^*_{ts} V_{td} (f_{ts} - f_{td} - f_{us} + f_{ud})]
\]

\[
+ V^*_{cb} V_{ub} [V^*_{cs} V_{cd} (f_{cb} - f_{cd} - f_{ub} + f_{ud}) + V^*_{ts} V_{td} (f_{tb} - f_{td} - f_{ub} + f_{ud})]
\]

with \(\lambda^U_{DD'} \equiv V^*_{UD'} V_{UD}, \lambda^D_{UU'} \equiv V^*_{UD} V_{U'D}, U = u, c, t \) and \(D = d, s, b, x_q = (m_q/m_W)^2\) and \(f_{UD} \equiv f(x_U, x_D)\)

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

Numerically, one obtains

\[
b_x \simeq 3.6 \cdot 10^{-7} e^{0.07 i}
\]

The quark masses are taken their values at \(m_c\) scale as given in \[4\]. The other contribution to the Lagrangian is the dipenguin and it gives \[57, 58, 62\].
\[
\Delta H = -\frac{G_F^2 \alpha_S}{8\pi^3} \left[ \lambda_{ca}^D E_0(x_D) \right] \left[ \lambda_{sd}^U E_0(x_U) \right] \bar{s} \gamma_\mu T^a d_L \left( g^{\mu\nu} \Box - \partial^\mu \partial^\nu \right) \bar{u} \gamma_\nu T^a c_L
\]

\[
= -\frac{G_F^2 \alpha_S}{8\pi^3} p_g \bar{s} \gamma_\mu T^a d_L \left( g^{\mu\nu} \Box - \partial^\mu \partial^\nu \right) \bar{u} \gamma_\nu T^a c_L
\]

\[
\equiv \frac{G_F^2 \alpha_S}{16\pi^3} P_g \mathcal{O}
\]

\[
p_g \equiv \left[ \lambda_{ca}^D E_0(x_D) \right] \left[ \lambda_{sd}^U E_0(x_U) \right] = \left[ V_{cs}^* V_{us} \left( E_0(x_s) - E_0(x_d) \right) + V_{cb}^* V_{ub} \left( E_0(x_b) - E_0(x_d) \right) \right] \left[ V_{cd} V_{cs} \left( E_0(x_c) - E_0(x_u) \right) + V_{td} V_{ts} \left( E_0(x_t) - E_0(x_u) \right) \right]
\]

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

\[
E_0(x) = \frac{1}{12(1-x)^4} \left[ x(1-x)(18-11x-x^2) - 2(4-16x+9x^2) \log(x) \right]
\]

The operator \( \mathcal{O} \) can be reduced as

\[
\mathcal{O} = \bar{s} \gamma_\mu T^a d_L \left( g^{\mu\nu} \Box - \partial^\mu \partial^\nu \right) \bar{u} \gamma_\nu T^a c_L = \bar{s} \gamma_\mu T^a d_L \Box \left( \bar{u} \gamma_\nu T^a c_L \right) + \bar{s} \partial T^a d_L \bar{u} \partial T^a c_L
\]

\[
= -q^2 \bar{s} \gamma_\mu T^a d_L \bar{u} \gamma_\nu T^a c_L - \left( m_s T^a d_{S-P} + m_d \bar{s} T^a d_{S+P} \right) \cdot \left( m_c \bar{u} T^a c_{S+P} + m_u \bar{u} T^a c_{S-P} \right) - q^2 \bar{s} \gamma_\mu T^a d_L \bar{u} \gamma_\nu T^a c_L - m_s m_c \bar{s} T^a d_L \bar{u} T^a c_R - m_d m_u \bar{s} T^a d_L \bar{u} T^a c_R
\]

where \( q^2 \) is the gluon momentum and \( N \) is the colour number. This expression can be
simplified using the fact that
\[
\hat{s} \gamma_\mu T^a d_L \bar{u} \gamma^\mu T^a c_L = \frac{1}{2} \left( O_1 - \frac{1}{N} O_2 \right)
\]
\[
\hat{s} T^a d_L \bar{u} T^a c_R = -\frac{1}{4} \hat{s} \gamma_\mu \gamma_\mu d_L - \frac{1}{2N} \hat{s} d_L \bar{u} c_R
\]
\[
\hat{s} T^a d_R \bar{u} T^a c_L = -\frac{1}{4} \hat{s} \gamma_\mu \gamma_\mu d_R - \frac{1}{2N} \hat{s} d_R \bar{u} c_L
\]
\[
\hat{s} T^a d_L \bar{u} T^a c_L = -\frac{1}{4} \hat{s} c_L \bar{u} d_L - \frac{1}{16} \hat{s} \sigma_\mu \nu c_L \bar{u} \sigma^{\mu \nu} d_L - \frac{1}{2N} \hat{s} d_L \bar{u} c_L
\]
\[
\hat{s} T^a d_R \bar{u} T^a c_R = -\frac{1}{4} \hat{s} c_R \bar{u} d_R - \frac{1}{16} \hat{s} \sigma_\mu \nu c_R \bar{u} \sigma^{\mu \nu} d_R - \frac{1}{2N} \hat{s} d_R \bar{u} c_R
\]  \hspace{1cm} (22)

Once taking the expectation values, one obtains
\[
\langle \mathcal{O} \rangle = -q^2 \langle \hat{s} \gamma_\mu T^a d_L \bar{u} \gamma^\mu T^a c_L \rangle - m_q m_c \langle \hat{s} T^a d_L \bar{u} T^a c_R \rangle - m_u m_u \langle \hat{s} T^a d_R \bar{u} T^a c_L \rangle
\]
\[
= -q^2 2 \left( 1 - \frac{1}{N^2} \right) X_{D^0 K^-}^+ + m_q m_c \left( 1 - \frac{1}{N} \right) X_{D^0 K^-}^+ + \frac{5m_d}{8N m_s} m_D^2 X_{K^- \pi^+}^0 \hspace{1cm} (23)
\]

Hence, one gets for the Wilson coefficients
\[
\Delta a_1 = -\frac{G_F m_w^2}{\sqrt{2} \pi^2 V_{cs}^* V_{ud} N} b_2 = \frac{G_F \alpha_S}{4\sqrt{2} \pi^3 V_{cs}^* V_{us}^*} \left[ \frac{q^2}{2} \left( 1 - \frac{1}{N^2} \right) - \frac{m_q m_c}{4} \left( 1 - \frac{1}{N} \right) \right] p_g
\]
\[
\approx 2.8 \times 10^{-8} e^{-0.004i}
\]
\[
\Delta a_2 = -\frac{G_F m_w^2}{\sqrt{2} \pi^2 V_{cs}^* V_{ud} N} b_2 = \frac{G_F \alpha_S}{4\sqrt{2} \pi^3 V_{cs}^* V_{us}^*} \frac{5m_d m_w^2}{8N m_s} p_g
\]
\[
\approx -2.0 \times 10^{-9} e^{0.07i} \hspace{1cm} (24)
\]

where to obtain the last result it has been used the fact that for the decay $D^0 \to K^- \pi^+$, one can approximate $q^2 = (p_c + p_u)^2 = (p_s + p_d)^2 \approx (p_D - p_\pi/2)^2 = (m_D^2 + m_K^2)/2 + 3m_\pi^2/4$, by assuming that $p_c \simeq p_D$ and $p_u \simeq p_\pi/2$ and $\alpha_S \simeq 0.3$. It should be noticed that the box contribution is dominated by the heavy quarks while the penguin is by the light ones. The direct CP asymmetry is then
\[
A_{CP} = \frac{|A|^2 - |\bar{A}|^2}{|A|^2 + |\bar{A}|^2} = \frac{2 |r| \sin(\phi_2 - \phi_1) \sin(\alpha_E))}{|1 + r|^2} = 1.4 \times 10^{-10} \hspace{1cm} (25)
\]
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 $\alpha_E$ 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 $\chi_{K^{-}\pi^+}^{D^0}$ and $\chi_{D^0K^-}^{+}$. This is important to take into account the FSI interactions as the first one is identified as E contribution and the second one is identified as T contribution. In the next subsections, we shall calculate the Wilson coefficient for different models of New Physics. The first case will be assuming to have 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 CP asymmetry coming from new charged Higgs-like scalar fields, applying to 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 exclude by precision data\[^{63–70}\]. Recent reviews on consequences of a fourth generation can be found in \[^{71–81}\].

The $B \to K\pi$ CP asymmetries puzzles is easily solved by a fourth generation \[^{82–84}\] with a mass within the following range\[^{82}\]:

$$400 \text{ GeV} < m_{u_4} < 600 \text{ GeV}. \quad (26)$$

The value of SM4 parameters compatibles with the high precision LEP measurements \[^{64–66, 69}\] are

$$m_{u_4} - m_{d_4} \simeq \left(1 + \frac{1}{3} \ln \frac{m_H}{115 \text{ GeV}}\right) \times 50 \text{ GeV} \quad (27)$$

$$|V_{ud_4}|, |V_{u_4d}| \lesssim 0.04 \quad (28)$$

where $V$ is the CKM quark mixing matrix which is now a $4 \times 4$ unitary matrix. The direct search limits from LEPII and CDF \[^{85–87}\] are given by:

$$m_{u_4} > 311 \text{ GeV} \quad (29)$$

$$m_{d_4} > 338 \text{ GeV}.$$
Direct search by Atlas and CMS coll. have excluded $m_{d_4} < 480$ GeV and $m_{q_4} < 350$ GeV [88–90], above the tree level unitarity limit, $m_{u_4} < \sqrt{4\pi/3} v \simeq 504$ GeV. But SM4 is far to be completely understood. Most of the experimental constraints are model-dependent. For instance it has been shown in [91] that the bound on $m_{u_4}$ should be relaxed up to $m_{u_4} > 350$ GeV if the decay $u_4 \rightarrow ht$ dominates. The recent LHC results which observe an excess in the $H \rightarrow \gamma\gamma$ corresponding to a Higgs mass around 125 GeV [92, 93] seems to exclude the SM4 scenario [94] but this results is based on the fact that once we include the next-to leading order electroweak corrections, the rate $\sigma(gg \rightarrow H) \times Br(H \rightarrow \gamma\gamma)$ is suppressed by more than 50% compared to the rate including only the leading order corrections [94–99]. This could be a signal of a non-perturbative 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 model with a fourth family is easy to compute as the contributions come from the same diagrams in the SM with just adding 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 $m_{t'} = 600$ GeV) to be

$$s_{14} = |V_{ub'}| = 0.017(14), \quad s_{24} = \frac{|V_{cb'}|}{c_{14}} = \frac{0.0084(62)}{c_{14}}, \quad s_{34} = \frac{|V_{tb'}|}{c_{14}c_{24}} = 0.07(8)$$

$$|V_{td}| = |V_{ts}| = 0.01(1), \quad |V_{ub}| = 0.07(8), \quad |V_{cb}| = 0.998(6), \quad |V_{tb}| \geq 0.98$$

$$\tan \theta_{12} = \frac{|V_{us}|}{|V_{ud}|}, \quad s_{13} = \frac{|V_{ub}|}{c_{14}}, \quad \delta_{13} = \gamma = 68^\circ$$

$$|V_{cb}| = |c_{13}c_{24}s_{23} - u_{13}^*u_{14}u_{24}| \simeq c_{13}c_{24}s_{23}$$

(30)

The two remaining phases ($\phi_{14}$ and $\phi_{24}$) 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

$$s_{13} = 0.00415, \quad s_{12} = 0.225, \quad s_{23} = 0.04, \quad s_{14} = 0.016, \quad s_{24} = 0.006, \quad s_{34} = 0.04$$

(31)

For a 4th sequential family the maxima value for the CP violation is obtained as

$$A_{CP} \simeq -1.1 \cdot 10^{-7}$$

(32)

where one uses $|V_{ub'}| = 0.06, \quad |V_{cb'}| = 0.03, \quad |V_{tb'}| = 0.25, \quad \phi_{14} = -2.9, \quad \phi_{24} = 1.3$
This maximal value is obtained when the parameters mentioned above are varied in a range allowed by the experiential constrains, according to eq. 30 in a ‘three sigma’ range. The phases are varied in the whole range, from \(-\pi\) to \(\pi\). Thus one can obtain an enhancement of thousand that may be large but still very far from the experimental possibilities.

B. A new charged gauge boson as 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 including a gauge \(SU(2)_R\) \[100–104\]. 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. \[105–109\]) and their parameters have been strongly constrained by experiments \[1, 110–114\]. Recently, CMS \[115, 116\] and ATLAS \[117, 118\] at LHC have improved the bound on the scale of the \(W_R\) gauge boson mass \[119\]. The new diagrams contributing to \(D \to K\pi\) are similar to the SM tree-level diagrams with \(W_L\) 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:

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

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 still it is suppressed due to the limit on \(M_{W_R}\) which has to be of order 2.3 TeV \[119\] in case of no-mixing Left right models.

In refs. \[120, 125\] 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 sector. This LR
muying is restricted by deviation to non-unitarity of the CKM quark mixing matrix. The results were that the Left-Right (LR) mixing angle called $\xi$ has to be smaller than 0.005\(^{121}\) and right scale $M_R$ bigger than 2.5 TeV\(^{119}\). If the Left-Right is not manifest (essentially that $g_R$ could be different from $g_L$ at Unification scale), the limit on $M_R$ scale is much less restrictive and the right gauge bosons could be as light as 0.3 TeV\(^{122}\). In such a case, $\xi$ can be as large as 0.02 if large CP violation phases in the right sector are present\(^{107}\) still compatible with experimental data\(^{123–125}\). Recently, precision measurement of the muon decay parameters done by TWIST collaboration\(^{126,127}\) put model independent limit on $\xi$ to be smaller than 0.03 (taking $g_L = g_R$). Let’s now compute the effect of the LR mixing gauge boson on our CP asymmetry. So first, one defines the charged current mixing matrix\(^{120}\)

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

(34)

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

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

(35)

where $V = V_{\text{CKM}}$ 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 of our process:

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

(36)

where $\alpha, \beta$ are color indices. It is easy to check that taking the limit $\xi \rightarrow 0$, one obtains eq.(5) with the only difference comes from the $c_2$ terms, the Fierz transformation has been applied. The terms of the effective Hamiltonian proportional to $\xi$ are:

$$
\Delta \mathcal{H}_{\text{eff}} \simeq \frac{G_F g_R}{\sqrt{2} g_L} \left[ c_1 \bar{s}_\gamma_\mu V^*_c L \bar{u}_R \gamma_\mu V_{ud}^R d_R + c_1 \bar{s}_\gamma_\mu V^*_c R \bar{u}_L \gamma_\mu V_{ud}^L d_L + c_2 \bar{s}_\alpha \gamma_\mu V^*_c L \bar{u}_R \gamma_\mu V_{ud}^R d_R + c_2 \bar{s}_\alpha \gamma_\mu V^*_c R \bar{u}_L \gamma_\mu V_{ud}^L d_L \right] + \text{h. c.}
$$

(37)
The contribution to the amplitude proportional to $\xi$ is then given by:

$$
\Delta A = -i G_F \frac{g_R}{\sqrt{2}} g_L \xi \left[ -c_1 V_{cs}^* V_{ud}^R \left( X_{D^0 K^-}^{\pi^+} + \frac{2}{N} X_{D^0 K^-}^{D^0} \right) + c_1 V_{cs}^* V_{ud} \left( X_{D^0 K^-}^{\pi^+} + \frac{2}{N} X_{D^0 K^-}^{D^0} \right) \right] 
$$

$$
- c_2 V_{cs}^* V_{ud}^R \left( 2 \chi^{D^0} X_{D^0 K^-}^{D^0} + \frac{1}{N} X_{D^0 K^-}^{\pi^+} \right) + c_2 V_{cs}^* V_{ud} \left( 2 \chi^{D^0} X_{D^0 K^-}^{D^0} + \frac{1}{N} X_{D^0 K^-}^{\pi^+} \right) \right] 
$$

$$
= \frac{i G_F g_R}{\sqrt{2} g_L} \xi \left( V_{cs}^* V_{ud} - V_{cs}^* V_{ud}^R \right) \left( a_1 X_{D^0 K^-}^{\pi^+} + 2 \chi^{D^0} a_2 X_{D^0 K^-}^{D^0} \right) 
$$

$$
= -\frac{g_R}{g_L} \xi \left( V_{cs}^* V_{ud} - V_{cs}^* V_{ud}^R \right) \left( T - 2 \chi^{D^0} E \right) \right][38]
$$

where $\chi^{\pi^+}$ and $\chi^{D^0}$ are defined as

$$
\chi^{\pi^+} = \frac{m_\pi^2}{(m_c - m_s)(m_u + m_d)} 
$$

$$
\chi^{D^0} = \frac{m_D^2}{(m_c + m_u)(m_s - m_d)} \right][39]
$$

The CP asymmetry becomes

$$
A_{CP} = \frac{4(g_R/g_L)\xi}{V_{cs}^* V_{ud}|1 + r|^2} \left( 1 + 2 \chi^{D^0} \right) \text{Im}(V_{cs}^* V_{ud} - V_{cs}^* V_{ud}^R) \text{Im}(r) \right][40]
$$

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 notice that to obtain this results, it has been used the fact that the chiralities don’t 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 quark in the loop are the down quarks contrarily to process like $b \rightarrow s \gamma$ where the quarks in the QCD corrections are the up quarks and in that case, a strong effect from top quarks could be expected [128–131]. 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 considering contribution to the effective Hamiltonian responsible of the $D^0 \rightarrow K^- \pi^+$ process due to a new charged Higgs fields. The simple SM extensions which include new charged Higgs fields are the two Higgs doublet models (2HDM) [132, 133]. Usually, it is used to classify these 2HDM in three types: type I, II or III (for a review see ref. [134]). In 2HDM type II models (like 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 \[135\]. 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 \[136\]:

$$m_{H^+} > 240 \text{GeV at } 95\% CL$$

(41)

2HDM type III is a general model where both Higgs couples to up and down quarks. Of course, this means that 2HDM type III can induce Flavor violation in Neutral Current and thus it can be used to strongly constrain the new parameters in the model. We shall focus our interest to the two Higgs doublet of type III as the other two can be obtained from type III taking some limits. In the 2HDM of type III, the Yukawa Lagrangian can be written as \[137, 138\]:

$$L_{\text{eff}}^Y = \bar{Q}_a^f L \left[ Y_{d}^{i} \epsilon_{ab} H_{d}^{b*} - \epsilon_{d}^{i} H_{u}^{a} \right] d_{i} R
$$

$$- \bar{Q}_a^f L \left[ Y_{u}^{i} \epsilon_{ab} H_{u}^{b} + \epsilon_{u}^{i} H_{d}^{a} \right] u_{i} R + \text{H.c.},$$

where $\epsilon_{ab}$ is the totally antisymmetric tensor, and $\epsilon_{ij}^d$ parametrizes the non-holomorphic corrections which couple up (down) quarks to the down (up) type Higgs doublet. After electroweak symmetry breaking, $L_{\text{eff}}^Y$ gives rise to the following charged Higss-quarks interaction Lagrangian:

$$L_{\text{eff}}^{H^{\pm}} = \bar{u}_f \Gamma_{u_{f}d_{i}}^{H^{\pm} LR_{\text{eff}}} P_R d_i + \bar{u}_f \Gamma_{u_{f}d_{i}}^{H^{\pm} RL_{\text{eff}}} P_L d_i ,$$

(43)

with \[138\]

$$\Gamma_{u_{f}d_{i}}^{H^{\pm} LR_{\text{eff}}} = \frac{3}{2} \sin \beta V_{fj} \left( \frac{m_{u_{i}}}{v_u} \delta_{ji} - \epsilon_{ji}^u \tan \beta \right),$$

$$\Gamma_{u_{f}d_{i}}^{H^{\pm} RL_{\text{eff}}} = \frac{3}{2} \cos \beta \left( \frac{m_{u_{i}}}{v_u} \delta_{ji} - \epsilon_{ji}^u \tan \beta \right) V_{ji}$$

(44)

Here $v_u$ and $v_d$ are the vacuum expectations 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 can compute the effective Hamiltonian resulting from the tree level exchanging charged Higgs diagram that governs the process under consideration namely,

$$H_{\text{eff}} = \frac{G_F}{\sqrt{2}} V_{c_{\beta}} V_{a_{\beta}} \sum_{i=1}^{4} C_{i}^{H^+}(\mu) Q_{i}^{H^+}(\mu),$$

(45)
where $C_i^H$ are the Wilson coefficients obtained by perturbative QCD running from $M_{H^\pm}$ scale to the scale $\mu$ 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

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

(46)

And the Wilson coefficients $C_i^H$, at the electroweak scale, are given by

$$
C_1^H = \frac{\sqrt{2}}{G_F V_{ud}^* V_{ub} m_W^2} \left( \sum_{j=1}^{3} \cos \beta V_{j1} \left( \frac{m_u}{v_u} \delta_{j1} - \epsilon_{j1}^{u*} \tan \beta \right) \right) \left( \sum_{k=1}^{3} \cos \beta V_{k2}^* \left( \frac{m_u}{v_u} \delta_{k2} - \epsilon_{k2}^{u*} \tan \beta \right) \right),
$$

$$
C_2^H = \frac{\sqrt{2}}{G_F V_{ud}^* V_{ub} m_W^2} \left( \sum_{j=1}^{3} \sin \beta V_{j1} \left( \frac{m_d}{v_d} \delta_{j1} - \epsilon_{j1}^{d*} \tan \beta \right) \right) \left( \sum_{k=1}^{3} \sin \beta V_{k2}^* \left( \frac{m_s}{v_s} \delta_{k2} - \epsilon_{k2}^{d*} \tan \beta \right) \right),
$$

$$
C_3^H = \frac{\sqrt{2}}{G_F V_{ud}^* V_{ub} m_W^2} \left( \sum_{j=1}^{3} \cos \beta V_{j1} \left( \frac{m_u}{v_u} \delta_{j1} - \epsilon_{j1}^{u*} \tan \beta \right) \right) \left( \sum_{k=1}^{3} \sin \beta V_{k2}^* \left( \frac{m_s}{v_s} \delta_{k2} - \epsilon_{k2}^{u*} \tan \beta \right) \right),
$$

$$
C_4^H = \frac{\sqrt{2}}{G_F V_{ud}^* V_{ub} m_W^2} \left( \sum_{k=1}^{3} \sin \beta V_{k2}^* \left( \frac{m_u}{v_u} \delta_{k2} - \epsilon_{k2}^{u*} \tan \beta \right) \right) \left( \sum_{j=1}^{3} \sin \beta V_{j1} \left( \frac{m_d}{v_d} \delta_{j1} - \epsilon_{j1}^{d*} \tan \beta \right) \right)
$$

(47)

We now discuss the experimental constraints on the $\epsilon_{ij}^q$ where $q = d, u$. The flavor-changing elements $\epsilon_{ij}^q$ for $i \neq j$ are strongly constrained from FCNC processes in the down sector because of tree-level neutral Higgs exchange. Thus, we are left with only $\epsilon_{11}^d, \epsilon_{22}^u$. Concerning the elements $\epsilon_{ij}^u$ we see that only $\epsilon_{11}^u, \epsilon_{22}^u$ can significantly affect the Wilson coefficients without any CKM suppression. Other $\epsilon_{ij}^u$ terms will be so small as the CKM suppression will be of orders $\lambda$ or $\lambda^2$ or higher and so we neglect them in our analysis. One of the important constraints that on $\epsilon_{ij}^q$ 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 \[138\]. 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 \[138\]

$$
|v_u(d)\epsilon_{ij}^{d(u)}| \leq |V_{ij}| \max \left[ m_{d(u_i)}, m_{d(u_j)} \right].
$$

(48)
which leads to

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

(49)

Clearly from the previous equation that \(\epsilon_{11}^u, \epsilon_{11}^d, \epsilon_{22}^d\) will be severely constrained by their small masses while \(\epsilon_{22}^u\) will be less constrained. Clearly from Eq.(49), the constraints imposed on \(\epsilon_{22}^u\) are tan \(\beta\) dependent. We now apply the constraints imposed on the real and imaginary parts of \(\epsilon_{22}^u\) 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 \(\epsilon_{22}^u\) thus they will be several order of magnitudes 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 \(\epsilon_{22}^u\) can be obtained from \(D - \bar{D}\) mixing, \(K - \bar{K}\) mixing. For \(K - \bar{K}\) mixing, the new contribution from charged Higgs mediation corresponding to top quark running in the loop will be much dominant than the contribution in the case where 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 \(\epsilon_{22}^u\) and \(\epsilon_{31}^u\) not on \(\epsilon_{32}^u\). In fact, as mentioned in ref.[138], the constraints on \(\epsilon_{32}^u\) and \(\epsilon_{31}^u\) are even weak and \(\epsilon_{32}^u\) and \(\epsilon_{31}^u\) can be sizeable. By a similar argument we can neither use the process \(b \rightarrow s\gamma\) nor the Electric dipole moment (EDM) to constraint \(\epsilon_{22}^u\).
Regarding $D - \bar{D}$ mixing one expects a similar situation like that in $K - \bar{K}$ about the dominance of 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

We take into accounts only box diagram that contribute to $D - \bar{D}$ mixing mediated by exchanging strange quark and charged Higgs. Other contributions from box diagram mediated by down or bottom quarks and charged Higgs are suppressed by the CKM factors. Since SM contribution to $D - \bar{D}$ mixing is very small we neglect its contribution and neglect its interference with charged Higgs mediation contribution. Thus effective Hamiltonian for this case can be written as:

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

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

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

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

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

$$C_3 = \frac{I_1(x_s)}{64\pi^2} \left( \sum_{j=1}^3 \sin \beta V_{2j}^* \left( \frac{m_u}{v_d} \delta_{j1} - \epsilon_{j1}^u \tan \beta \right) \right) \left( \sum_{k=1}^3 \sin \beta V_{1k} \left( \frac{m_s}{v_d} \delta_{k2} - \epsilon_{k2}^d \tan \beta \right) \right) \times \left( \sum_{l=1}^3 \cos \beta V_{1l} \left( \frac{m_u}{v_u} \delta_{l1} - \epsilon_{l1}^u \tan \beta \right) \right) \left( \sum_{n=1}^3 \cos \beta V_{2n} \left( \frac{m_c}{v_u} \delta_{n2} - \epsilon_{n2}^c \tan \beta \right) \right),$$

$$C_4 = \frac{m_s^2 I_2(x_s)}{16\pi^2 m_H^2} \left( \sum_{j=1}^3 \sin \beta V_{2j}^* \left( \frac{m_u}{v_d} \delta_{j1} - \epsilon_{j1}^u \tan \beta \right) \right) \left( \sum_{k=1}^3 \sin \beta V_{1k} \left( \frac{m_s}{v_d} \delta_{k2} - \epsilon_{k2}^d \tan \beta \right) \right) \times \left( \sum_{l=1}^3 \cos \beta V_{1l} \left( \frac{m_u}{v_u} \delta_{l1} - \epsilon_{l1}^u \tan \beta \right) \right) \left( \sum_{n=1}^3 \cos \beta V_{2n} \left( \frac{m_c}{v_u} \delta_{n2} - \epsilon_{n2}^c \tan \beta \right) \right).$$

(52)

where $x_s = m_s^2 / m_H^2$ and the integrals are defined as follows:

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

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

(53)

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

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

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

$$\tilde{C}_3 = C_3,$$

$$\tilde{C}_4 = C_4.$$  

(54)

Our set of operators $Q_1, Q_2$ and $Q_4$ given in Eq.(51) are equivalent to their corresponding operators given in Refs. [139, 140] 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. [139, 140] 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. [139, 140] since their matrix elements are equal. Thus our Wilson coefficients can be subjected to the constraints given in Ref. [139, 140] and so we find that

$$|C_1| \leq 5.7 \times 10^{-7} \left[ \frac{m_{H^\pm}}{1 \text{TeV}} \right]^2$$

$$|C_2| \leq 1.6 \times 10^{-7} \left[ \frac{m_{H^\pm}}{1 \text{TeV}} \right]^2$$

$$|C_3| \leq 3.2 \times 10^{-7} \left[ \frac{m_{H^\pm}}{1 \text{TeV}} \right]^2$$

$$|C_4| \leq 5.6 \times 10^{-8} \left[ \frac{m_{H^\pm}}{1 \text{TeV}} \right]^2$$

(55)

the constraints on $\tilde{C}_1 - \tilde{C}_4$ are similar to those $C_1 - C_4$. As can be seen from Eq. (55), the constraints on the Wilson coefficients will be strong for small charged Higgs masses. We can proceed now to derive the constraints on $\epsilon^{u}_{22}$ using the upper bound on $\tilde{C}_2$ for instance.

Keeping terms corresponding to first order in $\lambda$ where $\lambda$ is the CKM parameter we find that, for $m_{H^\pm} = 300$ GeV and $\tan \beta = 55$

$$\tilde{C}_2 \times 10^{12} \approx 3 \left( -53.6\epsilon^{u}_{12} - 12.7\epsilon^{d}_{22} + 0.007 \right)^2 \left( -12.4\epsilon^{u \ast}_{12} - 53.4\epsilon^{u \ast}_{22} + 0.007 \right)^2$$

(56)

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

$$\tilde{C}_2 \times 10^{14} \approx 3.6 \left( -487.1\epsilon^{d}_{12} - 115.0\epsilon^{d}_{22} + 0.06 \right)^2 \left( -112.5\epsilon^{u \ast}_{12} - 486.7\epsilon^{u \ast}_{22} + 0.007 \right)^2$$

(57)

In both Eqs. (56, 57) we can drop terms proportional to $\epsilon^{u \ast}_{12}$ to a good approximation as they have small coefficients in comparison to $\epsilon^{u}_{22}$ 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 $\epsilon^{d}_{12}$ can not be large to not allow flavor changing neutral currents and so we can drop terms proportional to $\epsilon^{d}_{12}$ in Eqs. (56, 57) to a good approximation also. thus we are left with $\epsilon^{d}_{22}$ and $\epsilon^{u}_{22}$ in both Eqs. (56, 57). Comparing their coefficients shows that $\epsilon^{u}_{22}$ has a large coefficient and thus we can drop $\epsilon^{d}_{22}$ terms. An alternative way is to assume that $\epsilon^{u}_{22}$ terms are the dominant ones in comparison to the other $\epsilon^{u,d}_{ij}$ terms and proceed to set upper bounds on $\epsilon^{u}_{22}$. 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 $\tilde{C}_2$, using Eq. (55), which reads in this case

$$|\tilde{C}_2| \leq 1.4 \times 10^{-8}$$

(58)
Clearly from Eqs. (56, 57, 58) the bounds that can be obtained on $\epsilon_{22}^u$ will be so loose and thus $D - D$ mixing can not lead to a strong constraints on $\epsilon_{22}^u$.

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 notations in Ref. [138], can be expressed as

$$B(D_q^+ \to \tau^+\nu) = \frac{G_F^2 |V_{cq}|^2}{8\pi} \frac{m_D^2 f_{D_q}^2 m_{D_q}}{m_\tau m_{D_q}} \left(1 - \frac{m_\tau^2}{m_{D_q}^2}\right)^2 \tau_{D_q} \times \left|1 + \frac{m_{D_q}^2}{(m_c + m_q) m_\tau} \frac{(C_{R}^{eq} - C_{L}^{eq})}{C_{SM}^{eq}}\right|^2.$$  \hfill (59)

Where we have used [141]

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

Where the SM Wilson coefficient is given by $C_{SM}^{eq} = 4 G_F V_{cq} / \sqrt{2}$ and the Wilson coefficients $C_{L}^{eq}$ and $C_{R}^{eq}$ at the matching scale are given by

$$C_{R(L)}^{eq} = \frac{-1}{M_{H^\pm}} \Gamma_{cq}^{LR(RL),H^\pm} \frac{m_c}{v} \tan \beta,$$  \hfill (61)

with the vacuum expectation value $v \approx 174$ GeV and $\Gamma_{cq}^{LR(RL),H^\pm}$ can be read from Eq. [14]. Setting the charged Higgs contribution to zero and $f_{D_q} = 248 \pm 2.5$ MeV [142], we find that $B^{SM}(D_d^+ \to \tau^+\nu) \approx 9.5 \times 10^{-4}$ and $B^{SM}(D^+ \to \tau^+\nu) = (5.11 \pm 0.11) \times 10^{-2}$ which is in close agreement with the results in Ref. [143–145]. The experimental values of these Branching ratios are given by $B(D_d^+ \to \tau^+\nu) < 2.1 \times 10^{-3}$ [146] while $B(D_s^+ \to \tau^+\nu) = (5.38 \pm 0.32) \times 10^{-2}$ [147]. Keeping the terms that are proportional to the dominant CKM elements we find for $q = d$

$$\Gamma_{cd}^{H^\pm LR eff} = \cos \beta V_{11} \left(-\epsilon_{12}^{u*} \tan \beta\right)$$

$$\Gamma_{cd}^{H^\pm LR eff} = \sin \beta V_{11} \left(\frac{m_d}{v_d} - \epsilon_{11}^{d} \tan \beta\right)$$  \hfill (62)
While for \( q = s \) we find

\[
\begin{align*}
\Gamma_{cs}^{H^\pm RL_{eff}} &= \cos \beta V_{22} \left( \frac{m_c}{v_u} - \epsilon_{22}^u \tan \beta \right) \\
\Gamma_{cs}^{H^\pm LR_{eff}} &= \sin \beta V_{22} \left( \frac{m_s}{v_d} - \epsilon_{22}^d \tan \beta \right)
\end{align*}
\]

(63)

Clearly from the last two equations, we need to consider the decay mode \( D_s^+ \rightarrow \tau^+ \nu \) to constrain \( \epsilon_{22}^u \). For \( \tan \beta = 10 \) we find that

\[
\begin{align*}
\Gamma_{cs}^{H^\pm RL_{eff}} \times 10^{-3} &\approx 0.71 - 968.6 \epsilon_{22}^u \\
\Gamma_{cs}^{H^\pm LR_{eff}} \times 10^{-3} &\approx 5.3 - 9686.0 \epsilon_{22}^d
\end{align*}
\]

(64)

Clearly the coefficient of \( \epsilon_{22}^d \) is one order of magnitude larger than \( \epsilon_{22}^u \) and for larger \( \tan \beta \) one expects to be larger than. However, \( \epsilon_{22}^d \) is severely constraint by naturalness criterion and thus we expect that the term proportional to \( \epsilon_{22}^u \) to be larger and thus in our analysis we can drop \( \epsilon_{22}^d \) term and proceed to obtain the required constraints. We show in Figs.\( \ref{fig:4} \)

\[ \begin{array}{c}
\text{FIG. 4. Constraints on } \epsilon_{22}^u \text{ from } B(D_s^+ \rightarrow \tau^+ \nu). \text{ Left plot corresponding to } \tan \beta = 200 \text{ while right plot corresponding to } \tan \beta = 500. \text{ In both cases we take } m_{H^\pm} = 200 \text{ GeV.} \\
\end{array} \]

the allowed regions for the real and imaginary parts of \( \epsilon_{22}^u \) corresponding to two different values of the charged Higgs mass namely, \( m_{H^\pm} = 200 \) and \( m_{H^\pm} = 300 \) 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 \( \tan \beta = 500 \), the constraints become loose with the increasing of \( m_{H^\pm} \). This is expected as Wilson coefficients of the charged Higgs are
FIG. 5. Constraints on $\epsilon_{22}^u$ from $B(D_s^+ \to \tau^+ \nu)$. Left plot corresponding to $\tan \beta = 350$ while right plot corresponding to $\tan \beta = 500$. In both cases we take $m_{H^\pm} = 300$ GeV.

inversely proportional to the square of $m_{H^\pm}$ and thus their contributions to $B(D_s^+ \to \tau^+ \nu)$ become small for large $m_{H^\pm}$ which in turn make the constraints obtained are loose. Another remark from the figure is that the constraints become strong with the increasing of the value of $\tan \beta$ which is expected also from Eq. (61). This in contrast to the constraints derived by applying the naturalness criterion where we showed that the constraints become loose with the increasing of the value of $\tan \beta$.

3. CP violation in Charged Higgs

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

$$A = \left(C_1^{SM} + \frac{1}{N} C_2^{SM} + \chi^{\pi^+} (C_1^H - C_4^H)\right) X_{D^0 K^-}^{\pi^+} - \left(C_2^{SM} + \frac{1}{N} C_1^{SM} + \frac{1}{2N} (C_1^H - C_4^H \chi^{D^0} C_4^H)\right) X_{K^- \pi^+}^{D^0}$$

with $X_{P_2P_3}^{P_1} = i f_{P_1} \Delta_{P_2P_3}^2 F_0^{P_2P_3}(m_{P_1}^2)$, $\Delta_{P_2P_3}^2 = m_{P_2}^2 - m_{P_3}^2$ and $\chi^{\pi^+}$ and $\chi^{D^0}$ are previously defined as

$$\chi^{\pi^+} = \frac{m_\pi^2}{(m_c - m_s)(m_u + m_d)}$$

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

The form of the amplitude, $A$, shows how charged Higgs contribution can affect only the short physics (Wilson coefficients) without any new effect on the long range physics (hadronic
parameters). Thus 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:

$$A = V_{cs}^* V_{ud} (T^{SM+H} + E^{SM+H})$$  \hspace{1cm} (67)$$

where

$$T^{SM+H} = 3.14 \times 10^{-6} \simeq \frac{G_F}{\sqrt{2}} a_1^{SM+H} f_\pi (m_D^2 - m_K^2) F^{DK}(m_\pi^2)$$

$$E^{SM+H} = 1.53 \times 10^{-6} e^{122\phi} \simeq \frac{G_F}{\sqrt{2}} a_2^{SM+H} f_D (m_K^2 - m_\pi^2) F^{K\pi}(m_D^2)$$  \hspace{1cm} (68)$$

where

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

$$= \left( a_1 + \Delta a_1 + \chi^+ (C_1^H - C_4^H) \right)$$  \hspace{1cm} (69)$$

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

The CP asymmetry can be obtained using the relation

$$A_{CP} = \frac{|A|^2 - |\bar{A}|^2}{|A|^2 + |\bar{A}|^2} = \frac{2|T^{SM+H}| |E^{SM+H}| \sin(\phi_1 - \phi_2) \sin(-\alpha_E)}{|T^{SM+H} + E^{SM+H}|^2}$$  \hspace{1cm} (71)$$

with $\phi_i = Arg[a_i^{SM+H}]$ and $\alpha_E = Arg(\chi_E)$. As an example let us take $Re(\epsilon_{22}^u) = 0.04$, $Im(\epsilon_{22}^u) = 0.03$ which is 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 $Re(\epsilon_{22}^u) = -0.1$, $Im(\epsilon_{22}^u) = -0.3$ which is allowed point for $\tan \beta = 500$ and $m_{H^\pm} = 300$ GeV. Repeating the same steps as above we find that $A_{CP} \simeq 5.3 \times 10^{-2}$. Clearly in charged Higgs models the predicted CP asymmetry is so sensitive to the value of $\tan \beta$ and to the value of Higgs mass.

V. CONCLUSION

In this paper, we have studied the Cabibbo favored non-leptonic $D^0$ 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 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 SM prediction of the CP asymmetry but still the predicted CP asymmetry is far of the reach of LHCb or SuperB factory as SuperKEKB. The most promising models are no-manifest Left-Right extension 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 next generation of charm or B factory. The non-observation 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 to solve at the same time the puzzle coming from $B \to \tau \nu$ and give a large contribution to this CP asymmetry depending on the charged Higgs masses and couplings. A maximal value of 5% can be reached with a Higgs mass of 300 GeV and large $\tan \beta$.

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Appendix A: Operators and other definitions

We start by defining $X_{P_2P_3}^{P_1}$, where $P_1$ denotes a pseudoscalar meson, as follows

$$X_{P_2P_3}^{P_1} = i f_{P_1} \Delta_{P_2P_3}^{P_1} F_0^{P_2P_3}(m_{P_1}^2)$$ (A1)
where $\Delta_{P_2P_3}^2 = m_{P_2}^2 - m_{P_3}^2$. In terms of $X_{P_2P_3}^{P_1}$ we find that

\[
\begin{align*}
&\langle \pi^+ | \bar{u} \gamma_\mu \gamma_5 d | 0 \rangle < K^- | \bar{s} \gamma_\mu c | D^0 \rangle = -X_{\pi K^-}^{\pi^+} \\
&\langle K^- \pi^+ | \bar{s} \gamma_\mu d | 0 \rangle < 0 | \bar{u} \gamma_\mu \gamma_5 c | D^0 \rangle = X_{D^0 K^-}^{\pi^+} \\
&\langle \pi^+ | \bar{u} \gamma_5 d | 0 \rangle < K^- | \bar{s} c | D^0 \rangle = -\frac{m^2_\pi}{(m_c - m_s)(m_u + m_d)} X_{\pi K^-}^{\pi^+} = -\chi^{\pi^+} X_{D^0 K^-}^{\pi^+} \\
&\langle K^- \pi^+ | \bar{s} d | 0 \rangle < 0 | \bar{u} \gamma_5 c | D^0 \rangle = -\frac{m^2_D}{(m_c + m_u)(m_s - m_d)} X_{D^0 K^-}^{\pi^+} = -\chi^{D^0} X_{K^- \pi^+}^{D^0}
\end{align*}
\]

(A2)

Using Eq. (A2) we get

\[
\begin{align*}
&\langle K^- \pi^+ | 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 < K^- | \bar{s} \gamma_\mu c_L | D^0 \rangle + \frac{1}{N} \langle K^- \pi^+ | \bar{s} \gamma_\mu d_L | 0 \rangle < 0 | \bar{u} \gamma_\mu c_L | D^0 \rangle = X_{\pi K^-}^{\pi^+} - \frac{1}{N} \chi^{\pi^+} X_{D^0 K^-}^{\pi^+} \\
&\langle K^- \pi^+ | 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 < 0 | \bar{u} \gamma_\mu c_L | D^0 \rangle = -X_{D^0 K^-}^{\pi^+} + \frac{1}{N} \chi^{D^0} X_{K^- \pi^+}^{\pi^+} \\
&\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 < K^- | \bar{s} \gamma_\mu c_R | D^0 \rangle + \frac{1}{N} \langle K^- \pi^+ | \bar{s} \gamma_\mu d_R | 0 \rangle < 0 | \bar{u} \gamma_\mu c_R | D^0 \rangle = -\langle K^- \pi^+ | O_1 | D^0 \rangle < K^- | \bar{s} \gamma_\mu c_R | D^0 \rangle > - K^- \pi^+ | O_2 | D^0 > \\
&\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 < K^- | \bar{s} \gamma_\mu c_L | D^0 \rangle + \frac{2}{N} \langle K^- \pi^+ | \bar{s} \gamma_\mu d_L | 0 \rangle < 0 | \bar{u} \gamma_\mu c_L | D^0 \rangle = -X_{\pi K^-}^{\pi^+} - \frac{2}{N} \chi^{\pi^+} X_{D^0 K^-}^{\pi^+} \\
&\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 < 0 | \bar{u} \gamma_\mu c_L | D^0 \rangle + \frac{2}{N} \langle \pi^+ | \bar{u} d_{S+P} | 0 \rangle < K^- | \bar{s} c_{S-P} | D^0 \rangle = -X_{D^0 K^-}^{\pi^+} + \frac{2}{N} \chi^{D^0} X_{K^- \pi^+}^{\pi^+} \\
&\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 < K^- | \bar{s} \gamma_\mu c_R | D^0 \rangle + \frac{2}{N} \langle K^- \pi^+ | \bar{s} \gamma_\mu d_R | 0 \rangle < 0 | \bar{u} \gamma_\mu c_R | D^0 \rangle = X_{D^0 K^-}^{\pi^+} + \frac{2}{N} \chi^{D^0} X_{K^- \pi^+}^{\pi^+} \\
&\langle K^- \pi^+ | \bar{u} \gamma_\mu c_R \bar{s} \gamma_\mu d_L | D^0 \rangle = \langle \pi^+ | \bar{u} \gamma_\mu d_L | 0 \rangle < K^- | \bar{s} \gamma_\mu c_R | D^0 \rangle + \frac{2}{N} \langle K^- \pi^+ | \bar{s} \gamma_\mu d_L | 0 \rangle < 0 | \bar{u} \gamma_\mu c_R | D^0 \rangle = X_{D^0 K^-}^{\pi^+} + \frac{2}{N} \chi^{D^0} X_{K^- \pi^+}^{\pi^+} \\
&\langle K^- \pi^+ | \bar{u} \gamma_\mu c_R \bar{s} \gamma_\mu d_L | D^0 \rangle = \langle \pi^+ | \bar{u} \gamma_\mu d_L | 0 \rangle < K^- | \bar{s} \gamma_\mu c_R | D^0 \rangle + \frac{2}{N} \langle K^- \pi^+ | \bar{s} \gamma_\mu d_L | 0 \rangle < 0 | \bar{u} \gamma_\mu c_R | D^0 \rangle = X_{D^0 K^-}^{\pi^+} - \frac{2}{N} \chi^{\pi^+} X_{D^0 K^-}^{\pi^+}
\end{align*}
\]

(A3)
and for the scalar ones

\begin{align}
<K^{-\pi^+}|\bar{s}\bar{c}_L \bar{u}d_L|D^0> &= \chi^{\pi^+} X^{\pi^+}_{D^0 K^-} - \frac{1}{2N} \chi^{D^0} X^{D^0}_{K^- \pi^+} \\
<K^{-\pi^+}|\bar{u}\bar{c}_L \bar{s}d_L|D^0> &= \chi^{D^0} X^{D^0}_{K^- \pi^+} - \frac{1}{2N} \chi^{\pi^+} X^{\pi^+}_{D^0 K^-} \\
<K^{-\pi^+}|\bar{s}\bar{c}_R \bar{u}d_R|D^0> &= -\chi^{\pi^+} X^{\pi^+}_{D^0 K^-} + \frac{1}{2N} \chi^{D^0} X^{D^0}_{K^- \pi^+} \\
<K^{-\pi^+}|\bar{u}\bar{c}_R \bar{s}d_R|D^0> &= -\chi^{D^0} X^{D^0}_{K^- \pi^+} + \frac{1}{2N} \chi^{\pi^+} X^{\pi^+}_{D^0 K^-} \\
<K^{-\pi^+}|\bar{s}\bar{c}_L \bar{u}d_R|D^0> &= -\chi^{\pi^+} X^{\pi^+}_{D^0 K^-} + \frac{1}{2N} \chi^{D^0} X^{D^0}_{K^- \pi^+} \\
<K^{-\pi^+}|\bar{u}\bar{c}_L \bar{s}d_R|D^0> &= \chi^{D^0} X^{D^0}_{K^- \pi^+} - \frac{1}{2N} \chi^{\pi^+} X^{\pi^+}_{D^0 K^-} \\
<K^{-\pi^+}|\bar{s}\bar{c}_R \bar{u}d_L|D^0> &= -\chi^{\pi^+} X^{\pi^+}_{D^0 K^-} - \frac{1}{2N} \chi^{D^0} X^{D^0}_{K^- \pi^+} \\
<K^{-\pi^+}|\bar{u}\bar{c}_R \bar{s}d_L|D^0> &= \chi^{D^0} X^{D^0}_{K^- \pi^+} - \frac{1}{2N} \chi^{\pi^+} X^{\pi^+}_{D^0 K^-}
\end{align}

(A4)

where the Fierz’s ordering has been used

\begin{align}
(\bar{\psi}_1 \Psi_2)_L(\bar{\psi}_3 \Psi_4)_L = (\bar{\psi}_1 \Psi_4)_L(\bar{\psi}_3 \Psi_2)_L, & \quad (\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} = & \quad -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_3 (1 \pm \gamma_5)\sigma_{\mu\nu}\psi_2 \\
2(T_a)_{\alpha\beta}(T_a)_{\gamma\delta} = & \quad \delta_{\alpha\delta}\delta_{\beta\gamma} - \frac{1}{N} \delta_{\alpha\beta}\delta_{\gamma\delta}, \quad (T_a)_{\alpha\beta}(T_a)_{\gamma\delta} = \frac{N^2 - 1}{2NC}\delta_{\alpha\delta}\delta_{\beta\gamma} - \frac{1}{NC}(T^a)_{\alpha\delta}(T^a)_{\beta\gamma}
\end{align}

(A5)

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