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

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Outline

- Introduction
- CP asymmetry within SM
- Extra SM fermion family
- Left Right modes
- Models with charged Higgs
- Summary

Introduction

Despite the great success of the SM in describing and predicting of many phenomena and observables, there are a number of issues that the SM fails to address:

- Gauge coupling unification: The three scale running coupling constants $\alpha_1(\mu), \alpha_2(\mu)$ and $\alpha_3(\mu)$ corresponding to the gauge group $SU_C(3) \otimes SU_L(2) \otimes U_Y(1)$ do not unify at some very high energy scale.
- The SM does not account for the neutrino masses.
- The SM does not have a viable candidate for the cold dark matter of the universe.
- The strength of the charge conjugation-parity (CP) violation in the SM is not sufficient to account for the cosmological baryon asymmetry of the universe.

All these issues indicate that new physics beyond the SM should exist. One way to search for signals for new physics beyond SM is 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. As an example, the decay mode is $D^0 \to K^-\pi^+$ and the observable is the direct CP asymmetry as I will show below:

CP asymmetry within SM

At quark level this decay mode is generated via $c \to su\bar{d}$ transitions. In general the Hamiltonian describing $D^0 \to K^-\pi^+$ can be expressed as

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

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:

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

- The amplitude can be calculated using the effective hamiltonian via calculating the matrix elemnt $\langle K^-\pi^+|\mathcal{H}|\bar{D}^0\rangle$. One way to calculate this matrix element is the so called naive factorization:

(3) $$\langle 0|\bar{q}_2\gamma^\mu\gamma_5 q_1|\bar{M}\rangle = i f_M P^\mu, \qquad \langle \bar{M}|\bar{q}_1\gamma^\mu\gamma_5 q_2|0\rangle = -i f_M P^\mu$$

$$\langle \bar{M}'|\bar{q}_2\gamma^\mu q_1|\bar{M}\rangle = f_+^{MM'}(q^2)\left(P^\mu - \frac{P\cdot q}{q^2}q^\mu\right) + f_0^{MM'}(q^2)\frac{P\cdot q}{q^2}q^\mu \qquad (4)$$

with $P = P_M + P_{M'}$ and $q = P_M - P_{M'}$ and $f_0^{MM'}(q^2), f_+^{MM'}(q^2)$ are form factors.

- One of the disadvantages of the naive factorization is that no strong CP conserving phases are obtained and therefore no CP violation is predicted.
- In addition, it is well known that final state interaction effects (FSI) are very important in these channels.
- 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 seem to produce appropriate strong phases. 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. This is the so called quark diagram approach. Within this approach, the amplitude is decomposed into parts corresponding to generic quark diagrams.

For the process under consideration, the main contributions are the tree level quark contribution (T) and exchange quark diagrams (E).

$$A_{D^0 \to K^- \pi^+} \quad \equiv \quad V_{cs}^* V_{ud}(T+E) \tag{5}$$

with

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

where in NFA they can be approximately written as

$$T \simeq \frac{G_F}{\sqrt{2}} a_1 f_\pi (m_D^2 - m_K^2) F_0^{DK}(m_\pi^2) \tag{7}$$

$$E \simeq -\frac{G_F}{\sqrt{2}} a_2 f_D (m_K^2 - m_\pi^2) F_0^{K\pi}(m_D^2) \tag{8}$$

Here $a_1 \equiv c_1 + c_2/N_c$ and $a_2 \equiv c_2 - c_1/N_C$. At this tree level, there is no source for weak CP violating phase as can be seen from the real factor $V_{cs}^* V_{ud}$ in the amplitude above and thus direct CP asymmetry will be zero at this level. One need to go to loop level to have weak CP violating phase

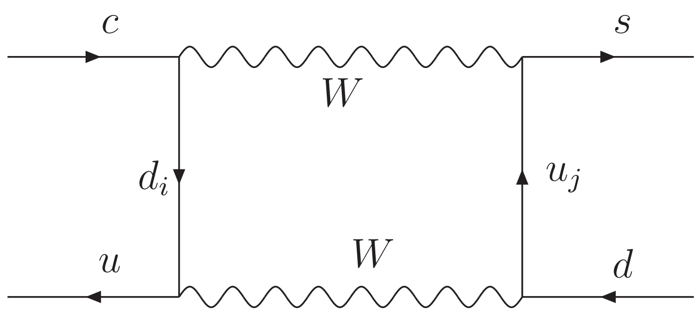

Figure 1: Box contribution.

$$\Delta\mathcal{H} \quad = \quad \frac{G_F^2 m_W^2}{2\pi^2} b_x \mathcal{O}_2$$

$$\begin{aligned} b_x \quad &= \quad V_{cs}^* V_{us}\left[V_{cs}^* V_{cd}\left(f_{cs} - f_{cd} - f_{us} + f_{ud}\right) + V_{ts}^* V_{td}\left(f_{ts} - f_{td} - f_{us} + f_{ud}\right)\right] \\ &\quad +V_{cb}^* V_{ub}\left[V_{cs}^* V_{cd}\left(f_{cb} - f_{cd} - f_{ub} + f_{ud}\right) + V_{ts}^* V_{td}\left(f_{tb} - f_{td} - f_{ub} + f_{ud}\right)\right] \end{aligned} \tag{9}$$

with $x_q = (m_q/m_W)^2$ and $f_{UD} \equiv f(x_U, x_D)$

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

One obtains $b_x \simeq 3.6 \cdot 10^{-7} \mathrm{e}^{0.07 \cdot i}$

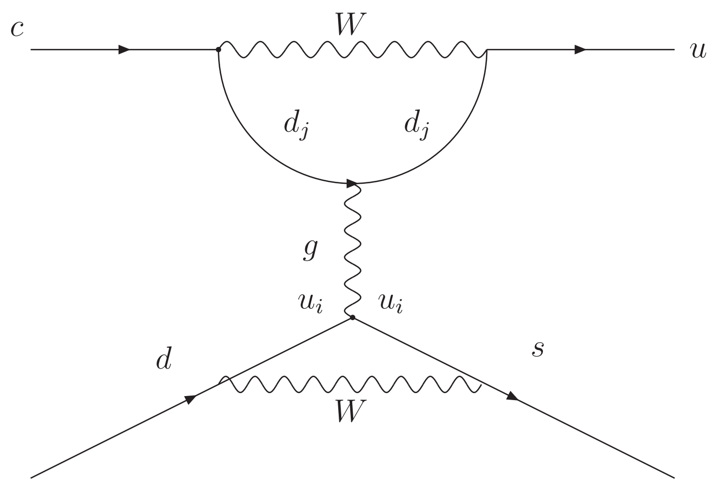

Figure 2: Di-penguins contribution.

Di-penguins contribution leads to two contributions to a_1 and a_2:

$$\begin{aligned}\Delta a_1 &= -\frac{G_F m_W^2}{\sqrt{2}\,\pi^2 V_{cs}^* V_{ud} N} b_x - \frac{G_F \alpha_S}{4\sqrt{2}\pi^3 V_{cs} V_{us}^*}\left[\frac{q^2}{2}\left(1-\frac{1}{N^2}\right) - \frac{m_c m_s}{4}\left(1-\frac{1}{N}\right)\right] p_g \\ &\simeq 2.8 \cdot 10^{-8} \mathrm{e}^{-0.004i}\end{aligned}$$

(10)

where N is the color factor and p_g is an expression contains the CKM factors and the loop integration in terms of Inami function.

$$\Delta a_2 = -\frac{G_F m_W^2}{\sqrt{2}\,\pi^2 V_{cs}^* V_{ud}} b_x - \frac{G_F \alpha_S}{4\sqrt{2}\pi^3 V_{cs} V_{us}^*} \frac{5 m_d m_D^2}{8 N m_s} p_g \simeq -2.0 \cdot 10^{-9} \mathrm{e}^{0.07i} \tag{11}$$

(12)

The direct CP asymmetry is then

$$A_{CP} = \frac{|A|^2 - |\bar{A}|^2}{|A|^2 + |\bar{A}|^2} = \frac{2|r|\sin(\phi_2 - \phi_1)\sin(\alpha_E)))}{|1+r|^2} = 1.4 \cdot 10^{-10} \tag{13}$$

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.(6).

Clearly within the SM the CP asymmetry is so tiny motivates investigation of new physics model beyond the SM. With New Physics, the general Hamiltonian is not only given by $\mathcal{O}_{1,2}$ and thus there will be new contributions from new Wilson coefficients with new sources for weak CP violating phases that may enhance the CP asymmetry significantly.

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 nor by the recent observation of the 125 Gev SM like Higgs boson. On the other hand most of the constraints on the masses of the new quarks are model dependent

- 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 $d_4 \equiv b'$.
- In Phys. Rev. D 83, 073008 (2011), it has been found that new CKM matrix elements can be obtained (all consistent with zero and for $m_{b'} = 600$ GeV) to be

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

- The two remaining phases (ϕ_{14} and ϕ_{24}) are unbounded an hence the absolute values of the CKM elements for the three families remain almost unchanged but not their phases.

- From these values one obtains

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

(15)

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

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

(16)

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 a the range allowed by the experimental constrains, according to eq. 14 in a 'three sigma' range.
- The phases are varied in the whole range, from $-\pi$ to π. Thus one can obtain an enhancement of thousand that may be large but still very far from the experimental possibilities.

Left Right models

As an example of such models, the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ defining the electroweak interaction.

- The new diagrams contributing to $D \to K\pi$ are similar to the SM tree-level diagrams with W_L is replaced by a W_R.
- Assuming no mixing between W_L and W_R gauge bosons we find:

$$\begin{aligned} \mathcal{H}_{\rm LR} &= \frac{G_F}{\sqrt{2}} \left(\frac{g_R m_W}{g_L m_{W_R}} \right)^2 V^*_{Rcs} V_{Rud} \left(c'_1 \bar{s}\gamma_\mu c_R \bar{u}\gamma^\mu d_R + c'_2 \bar{u}\gamma_\mu c_R \bar{s}\gamma^\mu d_R \right) + \mathrm{h.c.} \\ &= \frac{G_F}{\sqrt{2}} V^*_{cs} V_{ud} \left(c_1 \mathcal{O}_1 + c_2 \mathcal{O}_2 \right) + \mathrm{h.c.} \end{aligned} \quad (17)$$

  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.

- In this case the limits is $M_{W_R} \simeq 2.3$ TeV $\Longrightarrow$ the predicted CP asymmetry is still suppressed.

- Assuming mixing between W_L and W_R gauge bosons we find

$$\begin{pmatrix} W_L \\ W_R \end{pmatrix} = \begin{pmatrix} \cos\xi & -\sin\xi \\ \mathrm{e}^{i\omega}\sin\xi & \mathrm{e}^{i\omega}\cos\xi \end{pmatrix} \begin{pmatrix} W_1 \\ W_2 \end{pmatrix} \simeq \begin{pmatrix} 1 & -\xi \\ \mathrm{e}^{i\omega}\xi & \mathrm{e}^{i\omega} \end{pmatrix} \begin{pmatrix} W_1 \\ W_2 \end{pmatrix} \tag{18}$$

where W_1 and W_2 are the mass eigenstates. If the left-right symmetry is not manifest i.e. g_L is different than g_R at the unification scale, the limit on M_{W_R} is less restrictive and can be light as 0.3 TeV and $\xi \sim 10^{-2}$ if large CP violation phases in the right sector are present, Phys.Rev.D 40,1569 (1989).

- Assuming this scenario, one obtains the effective hamiltonian responsible of our process:

$$\begin{aligned} \mathcal{H}_{\text{eff.}} \quad &= \quad \frac{4G_F}{\sqrt{2}} \Bigg[c_1 \; \bar{s}\gamma_\mu \left(V^* P_L + \frac{g_R}{g_L} \xi \bar{V}^{R*} P_R \right)_{cs} c \;\; \bar{u}\gamma^\mu \left(V P_L + \frac{g_R}{g_L} \xi \bar{V} P_R \right)_{ud} d \\ &+ \quad c_2 \; \bar{s}_\alpha \gamma_\mu \left(V^* P_L + \frac{g_R}{g_L} \xi \bar{V}^{R*} P_R \right)_{cs} c_\beta \bar{u}_\beta \gamma^\mu \left(V P_L + \frac{g_R}{g_L} \xi \bar{V} P_R \right)_{ud} d_\alpha \Bigg] + \text{h. c.} \end{aligned} \tag{19}$$

where α, β are color indices.

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

$$\Delta A = -\frac{g_R}{g_L}\xi\left(\bar{V}_{cs}^{R*}V_{ud} - V_{cs}^{*}\bar{V}_{ud}^{R}\right)\left(T - 2\chi^{D^0}E\right) \quad (20)$$

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

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

$$\chi^{D^0} = \frac{m_D^2}{(m_c + m_u)(m_s - m_d)} \quad (21)$$

- The CP asymmetry becomes

$$A_{\mathrm{CP}} = \frac{4(g_R/g_L)\xi}{V_{cs}^{*}V_{ud}|1+r|^2}\left(1 + 2\chi^{D^0}\right)\mathrm{Im}\left(\bar{V}_{cs}^{R*}V_{ud} - V_{cs}^{*}\bar{V}_{ud}^{R}\right)\mathrm{Im}(r) \quad (22)$$

with $r = E/T$. For a value as large as $\xi \sim 10^{-2}$ the asymmetry can be as large as 0.1.

Models with Charged Higgs

We consider 2HDM of type III: the Yukawa Lagrangian can be written as :

(23)
$$\begin{aligned}\mathcal{L}_Y^{eff} &= \bar{Q}^a_{f\,L}\left[Y^d_{fi}\epsilon_{ab}H_d^{b\star} - \epsilon^d_{fi}H^a_u\right]d_{i\,R} \\ &- \bar{Q}^a_{f\,L}\left[Y^u_{fi}\epsilon_{ab}H_u^{b\star} + \epsilon^u_{fi}H^a_d\right]u_{i\,R} + \text{h.c.}\,,\end{aligned}$$

ϵ^q_{ij} parametrizes the non-holomorphic corrections which couple up (down) quarks to the down (up) type Higgs doublet.

- After electroweak symmetry breaking, $\mathcal{L}_Y^{eff}$ gives rise to the following charged Higss-quarks interaction Lagrangian:

(24)
$$\mathcal{L}^{eff}_{H^\pm} = \bar{u}_f \Gamma^{H^\pm\,LR\,\text{eff}}_{u_f d_i} P_R d_i + \bar{u}_f \Gamma^{H^\pm\,RL\,\text{eff}}_{u_f d_i} P_L d_i$$

with

$$\Gamma^{H^\pm\,LR\,\text{eff}}_{u_f d_i} = \sum_{j=1}^{3} \sin\beta\, V_{fj}\left(\frac{m_{d_i}}{v_d}\delta_{ji} - \epsilon^d_{ji}\tan\beta\right),$$

(25)
$$\Gamma^{H^\pm\,RL\,\text{eff}}_{u_f d_i} = \sum_{j=1}^{3} \cos\beta\left(\frac{m_{u_f}}{v_u}\delta_{jf} - \epsilon^{u\star}_{jf}\tan\beta\right)V_{ji}$$

this Lagrangian can be used to obtain the charged Higgs contribution to the effective Hamiltonian as

$$\mathcal{H}_{eff} = \frac{G_F}{\sqrt{2}} V_{cs}^* V_{ud} \sum_{i=1}^{4} C_i^H(\mu) Q_i^H(\mu), \tag{26}$$

where C_i^H are the Wilson coefficients and Q_i^H are the relevant local operators at low energy scale $\mu \simeq m_c$. The operators can be written as

$$\begin{aligned} Q_1^H &= (\bar{s}P_R c)(\bar{u}P_L d), \\ Q_2^H &= (\bar{s}P_L c)(\bar{u}P_R d), \\ Q_3^H &= (\bar{s}P_L c)(\bar{u}P_L d), \\ Q_4^H &= (\bar{s}P_R c)(\bar{u}P_R d), \end{aligned} \tag{27}$$

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

$$\begin{aligned} C_1^H &= \frac{\sqrt{2}}{G_F V_{cs}^* V_{ud} m_H^2} \left(\sum_{j=1}^{3} \cos\beta\, V_{j1} \left(\frac{m_u}{v_u} \delta_{j1} - \epsilon_{j1}^{u\star} \tan\beta \right) \right) \\ &\times \left(\sum_{k=1}^{3} \cos\beta\, V_{k2}^{\star} \left(\frac{m_c}{v_u} \delta_{k2} - \epsilon_{k2}^{u} \tan\beta \right) \right) \end{aligned} \tag{28}$$

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

(29)

In order to calculate the numerical values of the Wilson coefficients, we need to study the possible constraints on the parameters ϵ^{q}_{ij} where $q = d, u$ relevant to our decay mode.

- **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. 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

$$|v_{u(d)}\epsilon^{d(u)}_{ij}| \leq |V_{ij}| \ \max\left[m_{d_i(u_i)}, m_{d_j(u_j)}\right] \ . \quad (30)$$

  Clearly, $\epsilon^{u}_{11}, \epsilon^{d}_{11}, \epsilon^{d}_{22}$ will be severely constrained by their small masses while ϵ^{u}_{22} will be less constrained and these constraints will be $\tan\beta$ dependent.

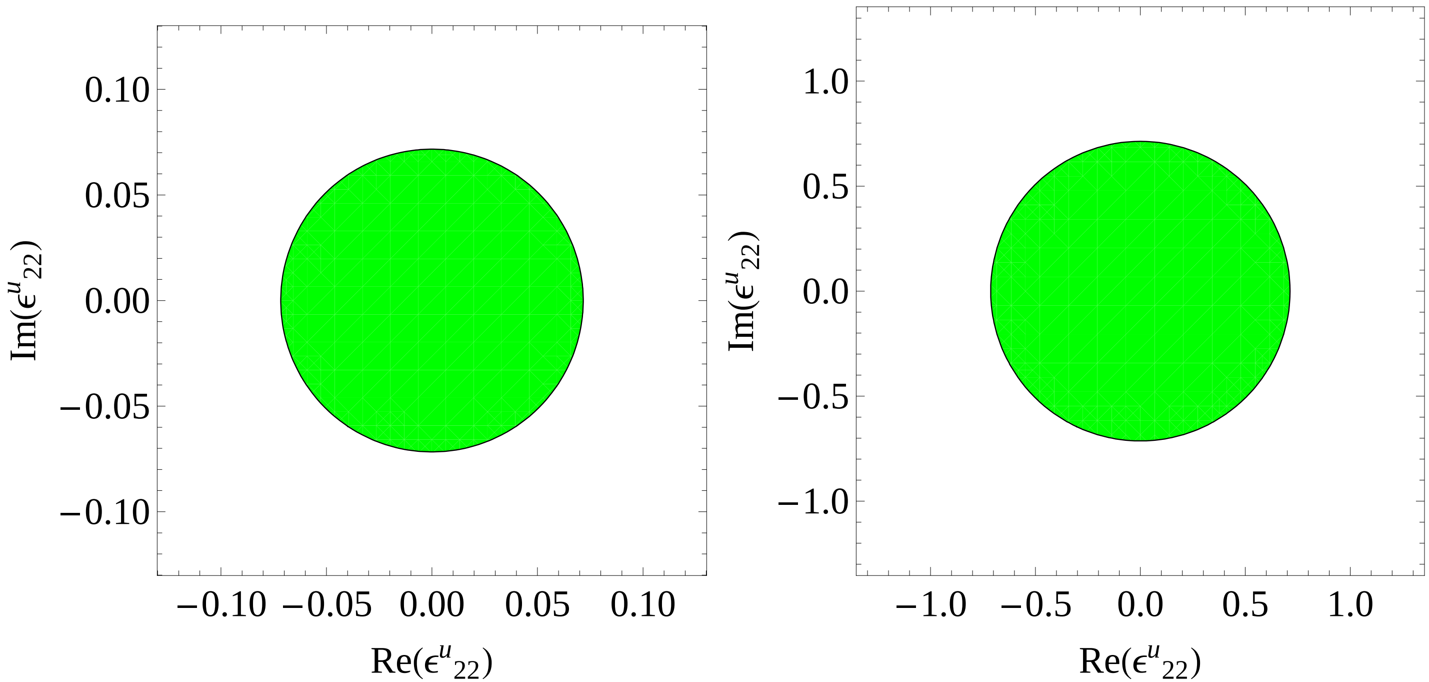

Figure 3: Constraints on ϵ^u_{22}. Left plot corresponding to $\tan\beta = 10$ while right plot corresponding to $\tan\beta = 100$.

- **$b \to s\gamma$, the Electric dipole moment and $K-\bar{K}$ mixing constraints**

  The constraints on ϵ^u_{22} from these processes are weak due to the dependency of the charged Higgs contributions on the ratio of the quark mass running in the loop to the charged Higgs mass. Thus the top quark contribution will be dominant than the charm quark contribution. As a consequence it is expected that constraints might be relevant on ϵ^u_{32} and ϵ^u_{31} not on ϵ^u_{22}.

- **$D-\bar{D}$ mixing constraints**

  Regarding $D-\bar{D}$ mixing one expects a similar situation like that in $K-\bar{K}$ mixing 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. Thus we need to consider possible constraints from $D-\bar{D}$ mixing. The charged Higgs mediation lead to the effective Hamiltonian:

$$\mathcal{H}^{|\Delta C|=2}_{H^\pm} = \frac{1}{m^2_{H^\pm}} \sum_{i=1}^{4} C_i(\mu) Q_i(\mu) + \tilde{C}_i(\mu)\tilde{Q}_i(\mu), \tag{31}$$

  where $C_i, \tilde{C}_i$ are the Wilson coefficients and $Q_i, \tilde{Q}_i$ are the relevant local operators at low energy scale

$$
\begin{aligned}
Q_1 &= (\bar{u}\gamma^\mu P_L c)(\bar{u}\gamma_\mu P_L c), \\
Q_2 &= (\bar{u} P_L c)(\bar{u} P_L c), \\
Q_3 &= (\bar{u}\gamma^\mu P_L c)(\bar{u}\gamma_\mu P_R c), \\
Q_4 &= (\bar{u} P_L c)(\bar{u} P_R c),
\end{aligned}
$$

(32)

$\tilde{Q}_i$ can be obtained from Q_i by changing the chirality $L \leftrightarrow R$.

- The Wilson coefficients C_i, are given by

$$
\begin{aligned}
C_1 &= \frac{I_1(x_s)}{64\pi^2}\left(\sum_{j=1}^{3} \sin\beta\, V_{2j}^* \left(\frac{m_s}{v_d}\delta_{j2} - \epsilon^d_{j2}\tan\beta\right)\right)^2 \left(\sum_{k=1}^{3} \sin\beta\, V_{1k} \left(\frac{m_s}{v_d}\delta_{k2} - \epsilon^d_{k2}\tan\right.\right. \\
C_2 &= \frac{m_s^2 I_2(x_s)}{16\pi^2 m_{H^\pm}^2}\left(\sum_{j=1}^{3} \sin\beta\, V_{2j}^* \left(\frac{m_s}{v_d}\delta_{j2} - \epsilon^d_{j2}\tan\beta\right)\right)^2 \\
&\times \left(\sum_{k=1}^{3} \cos\beta\, V_{k2} \left(\frac{m_u}{v_u}\delta_{k1} - \epsilon^{u\star}_{k1}\tan\beta\right)\right)^2,
\end{aligned}
$$

(33)

$$
\begin{aligned}
C_3 &= \frac{I_1(x_s)}{64\pi^2}\left(\sum_{j=1}^{3}\sin\beta\, V_{2j}^{*}\left(\frac{m_s}{v_d}\delta_{j2}-\epsilon_{j2}^{d}\tan\beta\right)\right)\left(\sum_{k=1}^{3}\sin\beta\, V_{1k}\left(\frac{m_s}{v_d}\delta_{k2}-\epsilon_{k2}^{d}\tan\beta\right)\right)\\
&\times \left(\sum_{l=1}^{3}\cos\beta\, V_{l2}\left(\frac{m_u}{v_u}\delta_{l1}-\epsilon_{l1}^{u\star}\tan\beta\right)\right)\left(\sum_{n=1}^{3}\cos\beta\, V_{n2}^{*}\left(\frac{m_c}{v_u}\delta_{n2}-\epsilon_{n2}^{u\star}\tan\beta\right)\right),\\
C_4 &= \frac{m_s^2 I_2(x_s)}{16\pi^2 m_{H^\pm}^2}\left(\sum_{j=1}^{3}\sin\beta\, V_{2j}^{*}\left(\frac{m_s}{v_d}\delta_{j2}-\epsilon_{j2}^{d}\tan\beta\right)\right)\left(\sum_{k=1}^{3}\sin\beta\, V_{1k}\left(\frac{m_s}{v_d}\delta_{k2}-\epsilon_{k2}^{d}\tan\beta\right)\right)\\
&\times \left(\sum_{l=1}^{3}\cos\beta\, V_{l2}\left(\frac{m_u}{v_u}\delta_{l1}-\epsilon_{l1}^{u\star}\tan\beta\right)\right)\left(\sum_{n=1}^{3}\cos\beta\, V_{n2}^{*}\left(\frac{m_c}{v_u}\delta_{n2}-\epsilon_{n2}^{u\star}\tan\beta\right)\right).
\end{aligned}
\tag{34}
$$

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

$$
\begin{aligned}
I_1(x_s) &= \frac{x_s+1}{(x_s-1)^2}+\frac{-2x_s\ln(x_s)}{(x_s-1)^3},\\
I_2(x_s) &= \frac{-2}{(x_s-1)^2}+\frac{(x_s+1)\ln(x_s)}{(x_s-1)^3}
\end{aligned}
\tag{35}
$$

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

$$\begin{aligned}
\tilde{C}_1 &= \frac{I_1(x_s)}{64\pi^2}\left(\sum_{j=1}^{3}\cos\beta\, V_{j2}\left(\frac{m_u}{v_u}\delta_{j1}-\epsilon^{u\star}_{j1}\tan\beta\right)\right)^2\left(\sum_{k=1}^{3}\cos\beta\, V^*_{k2}\left(\frac{m_c}{v_u}\delta_{k2}-\epsilon^{u\star}_{k2}\tan\beta\right.\right. \\
\tilde{C}_2 &= \frac{m_s^2 I_2(x_s)}{16\pi^2 m^2_{H^\pm}}\left(\sum_{j=1}^{3}\cos\beta\, V^*_{j2}\left(\frac{m_c}{v_u}\delta_{j2}-\epsilon^{u\star}_{j2}\tan\beta\right)\right)^2\left(\sum_{k=1}^{3}\sin\beta\, V_{1k}\left(\frac{m_s}{v_d}\delta_{k2}-\epsilon^{d}_{k2}\,\mathrm{ta}\right.\right. \\
\tilde{C}_3 &= C_3, \\
\tilde{C}_4 &= C_4.
\end{aligned}$$

(36)

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

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

(37)

Assuming that ϵ^{u}_{22} terms are the dominant ones in comparison to the other $\epsilon^{u,d}_{ij}$ terms and using the upper bound on $\tilde{C}_2$ from $D - \bar{D}$ mixing

$$| \tilde{C}_2 | \quad \leq \quad 1.4 \times 10^{-8} \tag{38}$$

we find that the bounds that can be obtained on ϵ^{u}_{22} will be so loos.

- **$D_s \to \tau\nu$ constraints**

  Within the 2HDM of type III, the charged Higgs can mediate these decay modes at tree level and hence the total branching ratios can be expressed as

$$\begin{aligned} \mathcal{B}(D_s^+ \to \tau^+\nu) &= \frac{G_F^2 |V_{cs}|^2}{8\pi} m_\tau^2 f_{D_s}^2 m_{D_s} \left(1 - \frac{m_\tau^2}{m_{D_s}^2}\right)^2 \tau_{D_s} \\ &\times \left| 1 + \frac{m_{D_s}^2}{(m_c + m_s)\, m_\tau} \frac{(C_R^{cs\,*} - C_L^{cs\,*})}{C_{SM}^{cs\,*}} \right|^2 . \end{aligned} \tag{39}$$

  Where $C_{\rm SM}^{cs} = 4 G_F\, V_{cs}/\sqrt{2}$ and

$$C^{cs}_{R(L)} = \frac{-1}{M_{H^\pm}^2}\, \Gamma_{cs}^{LR(RL),H^\pm}\, \frac{m_\tau}{v} \tan\beta \,, \tag{40}$$

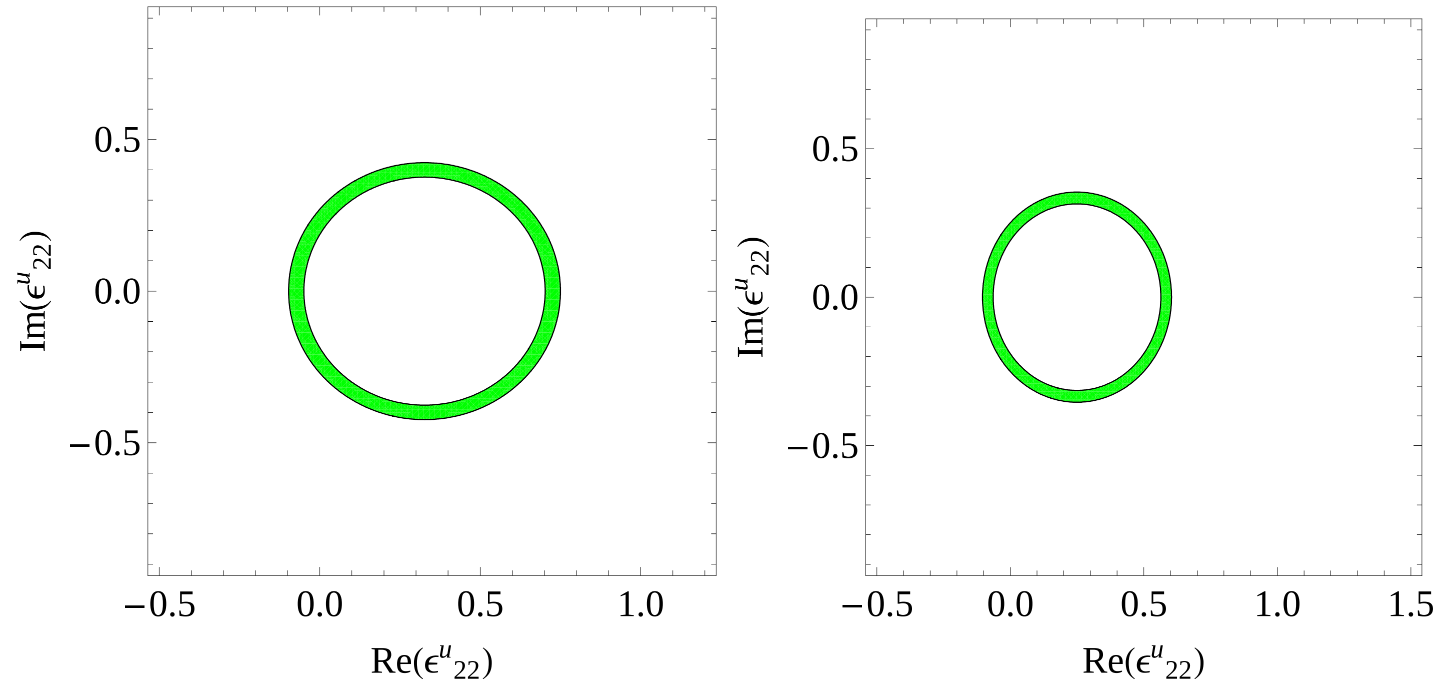

Figure 4: Allowed 2σ regions for the real and imaginary parts of ϵ^u_{22} from $\mathcal{B}(D_s^+ \to \tau^+\nu)$. Left plot corresponding to $\tan\beta = 117$ while right plot corresponding to $\tan\beta = 140$. In both cases we take $m_{H^\pm} = 150$ GeV.

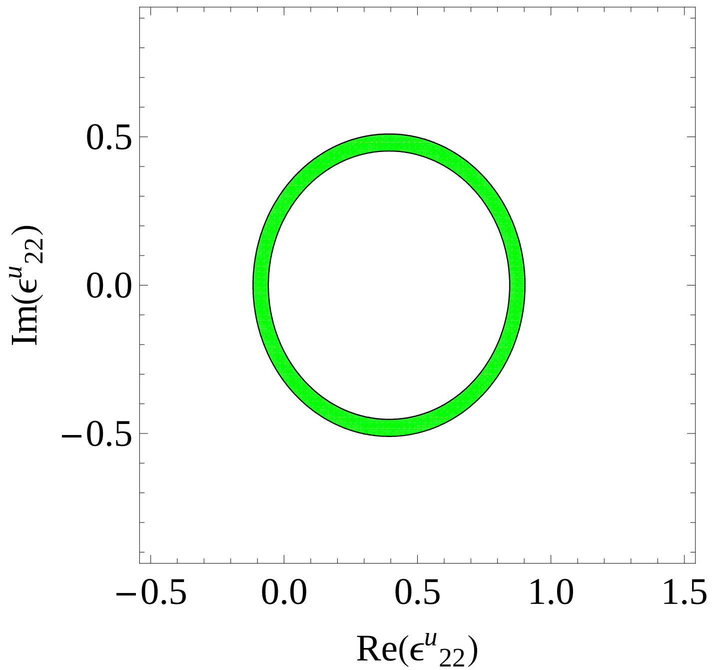

Figure 5: Allowed 2σ regions for the real and imaginary parts of ϵ^u_{22} from $\mathcal{B}(D_s^+ \to \tau^+\nu)$ corresponding to $\tan\beta = 140$ and $m_{H^\pm} = 180$ GeV.

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

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

(41)

with $X^{P_1}_{P_2 P_3} = i f_{P_1} \Delta^2_{P_2 P_3} F_0^{P_2 P_3}(m^2_{P_1})$, $\Delta^2_{P_2 P_3} = m^2_{P_2} - m^2_{P_3}$ and χ^{π^+} and χ^{D^0} are given as before. In terms of the amplitudes T and E introduced before in the case of the SM, we can write:

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

where

$$T^{SM+H} = 3.14 \times 10^{-6} \quad \simeq \quad \frac{G_F}{\sqrt{2}} a_1^{SM+H} f_\pi (m_D^2 - m_K^2) F_0^{DK}(m_\pi^2)$$

(43) $$E^{SM+H} = 1.53 \times 10^{-6} e^{122^\circ i} \quad \simeq \quad \frac{G_F}{\sqrt{2}} a_2^{SM+H} f_D (m_K^2 - m_\pi^2) F_0^{K\pi}(m_D^2)$$

where

(44) $$a_1^{SM+H} \quad = \quad \left(a_1 + \Delta a_1 + \chi^{\pi^+} (C_1^H - C_4^H)\right)$$

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

- The CP asymmetry can be obtained using the relation

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

with $\phi_i = Arg[a_i^{SM+H}]$ and α_E is the strong phase.

- As an example, taking $Re(\epsilon_{22}^u) = 0.04$, $Im(\epsilon_{22}^u) = 0.03$ which is allowed point for $\tan\beta = 10$. In this case we find that for a value of $m_{H^\pm} = 500$ GeV, $A_{CP} \simeq -3.7 \times 10^{-5}$ while for $m_H = 300$ GeV we find that $A_{CP} \simeq -1 \times 10^{-4}$.
- Another example where $Re(\epsilon_{22}^u) = -0.5$, $Im(\epsilon_{22}^u) = -0.3$ which is allowed point for $\tan\beta = 100$, 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 so sensitive to the value of $\tan\beta$ and to the value of Higgs mass.

Summary

- In this talk 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.
- 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 such as SuperKEKB.
- The most promising models are non-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 the 2HDM type III a maximal value of 1.5% can be reached with a Higgs mass of 250 GeV and large tanβ