Implications of the LHCb discovery of CP violation in charm decays

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Abstract

The recent measurement of ΔA_{CP} by the LHCb collaboration requires an $\mathcal{O}(10)$ enhancement coming from hadronic physics in order to be explained within the SM. We examine to what extent can NP models explain ΔA_{CP} without such enhancements. We discuss the implications in terms of a low energy effective theory as well as in the context of several explicit NP models.

1 Introduction to ΔA_{CP}

The LHCb experiment has announced discovery of direct CP violation in singly Cabibbo suppressed D decays [1],

$$\begin{aligned}\Delta A_{CP} &\equiv A_{CP}(K^+K^-) - A_{CP}(\pi^+\pi^-) \\ &= (-1.54 \pm 0.29) \times 10^{-3}.\end{aligned} \tag{1}$$

Here

$$A_{CP}(f) \equiv \frac{\Gamma(D^0 \to f) - \Gamma(\overline{D}^0 \to f)}{\Gamma(D^0 \to f) + \Gamma(\overline{D}^0 \to f)}. \tag{2}$$

In ΔA_{CP} effects of indirect CP violation approximately cancel out [2]. (Due to different decay time acceptances between the K^+K^- and $\pi^+\pi^-$ modes, a small residual effect of indirect CP violation remains.) Thus, ΔA_{CP} is a manifestation of CP violation in decay. The updated world average for the direct and indirect CP violating contributions to this asymmetry are [3]

$$\Delta A_{CP}^{\rm dir} = (-1.64 \pm 0.28) \times 10^{-3}, \tag{3}$$

$$A_{CP}^{\rm ind} = (+0.28 \pm 0.26) \times 10^{-3}. \tag{4}$$

The singly Cabibbo suppressed D^0 ($\overline{D}^0$) decay amplitudes A_f ($\overline{A}_f$) to a final CP eigenstate f can be written as [2]

$$\begin{aligned}A_f &= A_f^T e^{i\phi_f^T}\left[1 + r_f e^{i(\delta_f+\phi_f)}\right], \\ \overline{A}_f &= \eta_{CP} A_f^T e^{-i\phi_f^T}\left[1 + r_f e^{i(\delta_f-\phi_f)}\right],\end{aligned} \tag{5}$$

where $\eta_{CP} = \pm 1$ is the CP eigenvalue of f, the dominant singly Cabibbo suppressed "tree" amplitude is denoted by $A_f^T e^{\pm i\phi_f^T}$, and r_f parameterizes the relative magnitude of all subleading amplitudes (often called "penguin" amplitudes), which carry different strong (δ_f) and weak (ϕ_f) phases. Then

$$A_{CP}^{\rm dir}(f) = -\frac{2r_f \sin\delta_f \sin\phi_f}{1 + 2r_f \cos\delta_f \cos\phi_f + r_f^2}. \tag{6}$$

The Standard Model (SM) contribution to the individual asymmetries is CKM suppressed by a factor of

$$I_{\rm CKM} \equiv 2\mathcal{I}m\left(\frac{V_{ub}V_{cb}^*}{V_{us}V_{cs}^*}\right) \approx 1.4 \times 10^{-3}. \tag{7}$$

Naively, there is a further loop suppression by a factor of order $\alpha_s/\pi \sim 0.1$. One cannot exclude an enhancement factor of order 10 from hadronic physics [4–7], in which case (3) is accounted for by SM physics. Yet, it is not implausible that new physics (NP) dominates ΔA_{CP} [8, 9] (indeed, QCD-based LCSR calculations [10] support the latter option.)

In the following we assume that hadronic factors do not significantly alter the magnitude of the relevant effects; Thus, NP is required to explain the measured ΔA_{CP}. We analyze the implications of Eq. (3) on candidate models. We phrase our findings in terms of which NP models can or cannot account for the measurement, assuming that the SM contribution is negligible. Relaxing this assumption, the same implications can be conservatively read as upper bounds on the NP model parameters.

In 2011, experimental evidence for ΔA_{CP} [11] prompted several related studies [4, 5, 8, 12–14]. We provide an update to some of the relevant results, taking into account the recent discovery with a central value smaller by a factor of ~ 4 as well as all applicable existing bounds.

We begin with an effective field theory (EFT) analysis in Section 2. We follow with specific examples of models in which the measured ΔA_{CP} is explained: 2HDM in Section 3, the MSSM in Section 4 and models with vector-like up-quarks in Section 5. We conclude in Section 6.

2 Non-renormalizable operators

The relevant effects of new physics at a scale much higher than the electroweak breaking scale can be represented by the following effective Hamiltonian [8]:

$$\mathcal{H}^{\text{eff}-\text{NP}}_{|\Delta c|=1} = \frac{G_F}{\sqrt{2}} \sum_{i=1,2,5,6} \sum_q (C_i^q Q_i^q + C_i^{q\prime} Q_i^{q\prime}) + \frac{G_F}{\sqrt{2}} \sum_{i=7,8} (C_i^q Q_i + C_i^\prime Q_i^\prime) + \text{h.c.}, \tag{8}$$

where $q = \{d, s, b, u, c\}$, the list of operators includes

$$\begin{aligned}
Q_1^q &= (\bar{u}q)_{V-A}(\bar{q}c)_{V-A}, & Q_7 &= -\frac{e}{8\pi^2} m_c \bar{u}\sigma_{\mu\nu}(1+\gamma_5)F^{\mu\nu}c, \\
Q_2^q &= (\bar{u}_\alpha q_\beta)_{V-A}(\bar{q}_\beta c_\alpha)_{V-A}, & Q_8 &= -\frac{g_s}{8\pi^2} m_c \bar{u}\sigma_{\mu\nu}(1+\gamma_5)T^a G_a^{\mu\nu}c, \\
Q_5^q &= (\bar{u}c)_{V-A}(\bar{q}q)_{V+A}, \\
Q_6^q &= (\bar{u}_\alpha c_\beta)_{V-A}(\bar{q}_\beta q_\alpha)_{V+A},
\end{aligned} \tag{9}$$

and the primed operators are related to the non-primed ones via $A \leftrightarrow -A$ and $\gamma_5 \leftrightarrow -\gamma_5$.

The SM and NP contributions to ΔA_{CP} can be parameterized as

$$\Delta A_{CP} \approx I_{\text{CKM}} \frac{\alpha_s(m_c)}{\pi} \mathcal{I}m(\Delta R^{\text{SM}}) + \frac{2}{|V_{us}V_{cs}|} \sum_i \mathcal{I}m(C_i^{\text{NP}}) \mathcal{I}m(\Delta R_i^{\text{NP}}), \tag{10}$$

where $\Delta R^{\text{SM,NP}} = R_K^{\text{SM,NP}} + R_\pi^{\text{SM,NP}}$, and $R_K^{\text{SM,NP}}$ are the ratios of subleading amplitudes to the leading SM amplitude, after factorizing out the CKM dependence and the Wilson coefficient (the loop factor for R_K^{SM}). Thus the SM alone can explain the measured value of ΔA_{CP} for $\mathcal{I}m(\Delta R^{\text{SM}}) \approx 13$. In the following we conversely adopt the naive expectation, $\mathcal{I}m(\Delta R^{\text{SM}}) \sim \mathcal{I}m(\Delta R^{\text{NP}}) \sim 2$ (the factor of 2 is inspired by the U-spin limit, in which $A^{\text{SM}}_{K^+K^-} \approx -A^{\text{SM}}_{\pi^+\pi^-}$.) With this assumption, the measurement requires the existence of NP with a Wilson coefficient satisfying

$$\mathcal{I}m(C_i^{\text{NP}}) \sim \frac{\Delta A_{CP}}{18.2} \sim 9 \times 10^{-5}, \tag{11}$$

and the scale of NP can naively be estimated as $\lesssim 37$ TeV.

2.1 Constraints from $D^0 - \overline{D}^0$ mixing

The Hamiltonian of Eq. (8) is related to the effective Hamiltonian relevant for $|\Delta c| = 2$ transitions,

$$\mathcal{H}^{\text{eff}}_{|\Delta c|=2} = \frac{G_F}{\sqrt{2}} \left(\sum_{i=1}^{5} C_i^{cu} Q_i^{cu} + \sum_{i=1}^{3} C_i^{cu\prime} Q_i^{cu\prime} \right), \tag{12}$$

Table 1: Upper bounds on CP violating $\Delta c = 1$ operators from $D^0 - \overline{D}^0$ mixing, at the hadronic charm scale $\mu \approx 2\,\text{GeV}$.

f	$s-d$	$8d$
$\mathcal{I}m(C_{1,2}^{(f)})$	3.6×10^{-7}	9.6×10^{-4}
$\mathcal{I}m(C_5^{(f)\prime})$	5.6×10^{-8}	1.5×10^{-4}
$\mathcal{I}m(C_6^{(f)\prime})$	2.0×10^{-8}	5.3×10^{-5}

where

$$\begin{aligned} Q_1^{cu} &= (\bar{u}c)_{V-A}(\bar{u}c)_{V-A} \qquad & Q_2^{cu} = (\bar{u}c)_{S-P}(\bar{u}c)_{S-P} \\ Q_3^{cu} &= (\bar{u}_\alpha c_\beta)_{S-P}(\bar{u}_\beta c_\alpha)_{S-P} \qquad & Q_4^{cu} = (\bar{u}c)_{S-P}(\bar{u}c)_{S+P} \\ Q_5^{cu} &= (\bar{u}_\alpha c_\beta)_{S-P}(\bar{u}_\beta c_\alpha)_{S+P}. & \end{aligned} \tag{13}$$

The contributions of $\mathcal{H}^{\text{eff}}_{|\Delta c|=2}$ to $D^0 - \overline{D}^0$ mixing are computed using the following formula:

$$\langle \bar{D}^0 | \mathcal{H}^{\text{eff}}_{|\Delta c|=2} | D^0 \rangle_i = \frac{G_F}{\sqrt{2}} \sum_{j=1}^{5} \sum_{r=1}^{5} \left(b_j^{(r,i)} + \eta c_j^{(r,i)} \right) \eta^{a_j} \times C_i^{cu}(\mu) \langle \bar{D}^0 | Q_r^{cu} | D^0 \rangle, \tag{14}$$

where all relevant parameters and hadronic matrix elements are defined in Ref. [15].

Using the up-to-date 95% C.L regions for the mixing parameters [3],

$$\begin{aligned} x_{12} &\in [0.22, 0.63]\% \\ y_{12} &\in [0.59, 0.75]\% \\ \phi_{12} &\in [-2.5^o, 1.8^o], \end{aligned} \tag{15}$$

we obtain the following bounds:

$$\begin{aligned} \text{Im}(C_1^{cu}) &\lesssim 1.6 \times 10^{-9}; \qquad & \text{Re}(C_1^{cu}) \lesssim 3.6 \times 10^{-8}, \\ \text{Im}(C_4^{cu}) &\lesssim 1.7 \times 10^{-10}; \qquad & \text{Re}(C_4^{cu}) \lesssim 4.0 \times 10^{-9}, \\ \text{Im}(C_5^{cu}) &\lesssim 4.9 \times 10^{-10}; \qquad & \text{Re}(C_5^{cu}) \lesssim 1.1 \times 10^{-8}. \end{aligned} \tag{16}$$

Following Ref. [8], we can relate the two sets of Wilson coefficients via

$$\begin{aligned} C_1^{cu} &= \delta C_1^{cu} + \frac{g^2}{32\pi^2} \sum_q \lambda_q (C_2^q - C_1^q) \ln \frac{\mu^2}{m_W^2}, \\ C_4^{cu} &= \delta C_4^{cu} - \frac{g^2}{16\pi^2} \sum_q \lambda_q C_6^{q\prime} \ln \frac{\mu^2}{m_W^2}, \\ C_5^{cu} &= \delta C_5^{cu} - \frac{g^2}{16\pi^2} \sum_q \lambda_q C_5^{q\prime} \ln \frac{\mu^2}{m_W^2}. \end{aligned} \tag{17}$$

We then change basis to $Q_i^{s-d} = Q_i^s - Q_i^d, Q_i^{8d} = Q_i^s + Q_i^d - 2Q_i^b$, and take the counter-terms to zero to arrive at the bounds on the $\Delta c = 1$ operators, presented in Table 1. We conclude that the operators $Q_{1,2}^{(s-d)}$, $Q_{5,6}^{(s-d)\prime}$ and $Q_6^{(8d)\prime}$ cannot account for ΔA_{CP}.

2.2 Constraints from ϵ'/ϵ

Following Ref. [8], we use the master formula for ϵ'/ϵ, evaluating the matrix elements induced by the $|\Delta s| = 1$ operators at the large N_c limit. The NP contribution is then given by

$$\left|\frac{\epsilon'}{\epsilon}\right|_{\rm NP} \approx 10^2\Big|\mathcal{I}m\Big[3.5C_1^{(3/2)} + 3.4C_2^{(3/2)} - 1.7\rho^2 C_5^{(3/2)} - 5.2\rho^2 C_6^{(3/2)} - 0.04C_1^{(1/2)} - 0.12C_2^{(1/2)} - 0.04\rho^2 C_5^{(1/2)} + 0.11\rho^2 C_6^{(1/2)}\Big]\Big|, \tag{18}$$

where $C_i^{(3/2)} = \frac{1}{2}(-C_i^{(s-d)} + C_i^{(c-u)} + C_i^{(8d)}) + \frac{5}{4}C_i^{(b)}$, $C_i^{(1/2)} = \frac{1}{2}(C_i^{(s-d)} + C_i^{(c-u)} - C_i^{(8d)}) + \frac{1}{4}C_i^{(b)} - C_i^{(0)}$, and $\rho = m_K/m_s$. Taking the conservative bound $|\epsilon'/\epsilon|_{\rm NP} < |\epsilon'/\epsilon|_{\rm exp} \approx 1.7\times 10^{-3}$, the imaginary parts of the $|\Delta s| = 1$ Wilson coefficients are constrained. These are related to the $|\Delta c| = 1$ coefficients of interest via

$$C_i^{q(ds)} = \delta C_i^{q(ds)} + C_i^q \frac{g^2}{32\pi^2}\ln\frac{\mu^2}{m_W^2}. \tag{19}$$

The resulting bounds on the $|\Delta c| = 1$ Wilson coefficients are presented in Table 2. Comparing these bounds to Eq. (11), we conclude that the operators $Q_{5,6}^{(f)}$ with $f \in \{s-d, c-u, 8d, b\}$ cannot account for ΔA_{CP}.

Table 2: Upper bounds on CP violating $\Delta c = 1$ operators from $|\epsilon'/\epsilon|$, at the hadronic charm scale $\mu \approx 2\,\mathrm{GeV}$.

f	$s-d$	$c-u$	$8d$	b	0
$\mathcal{I}m(C_1^{(f)})$	4.8×10^{-4}	4.9×10^{-4}	4.8×10^{-4}	1.9×10^{-4}	2.1×10^{-2}
$\mathcal{I}m(C_2^{(f)})$	4.8×10^{-4}	5.0×10^{-4}	4.8×10^{-4}	2.0×10^{-4}	6.9×10^{-3}
$\mathcal{I}m(C_5^{(f)})$	3.6×10^{-5}	3.4×10^{-5}	3.6×10^{-5}	1.4×10^{-5}	7.4×10^{-4}
$\mathcal{I}m(C_6^{(f)})$	1.1×10^{-5}	1.1×10^{-5}	1.1×10^{-5}	4.6×10^{-6}	2.7×10^{-4}

We note that the set of operators, $\{Q_{7,8}, Q'_{7,8}, \forall f\, Q_{1,2}^{f\prime}, Q_{5,6}^{(c-u,b,0)\prime}\}$, are relevant to neither $D^0-\overline{D}^0$ mixing nor $|\epsilon'/\epsilon|$, and therefore are unconstrained. Table 3 summarizes which $\Delta c = 1$ operators can contribute to ΔA_{CP} at a level comparable to the current measured value.

Table 3: Classification of new physics operators Q_i according to whether upper bounds on $\mathcal{I}m(C_i^{\rm NP})$ from $D^0-\overline{D}^0$ mixing and ϵ'/ϵ are (i) much weaker than 9×10^{-5} ("allowed"), (ii) of order 9×10^{-5} ("marginal"), or (iii) much stronger than 9×10^{-5} ("disfavored").

Allowed	Marginal	Disfavored
$Q_{7,8}$, $Q'_{7,8}$,	$Q_5^{(8d)\prime}$	$Q_{1,2}^{(s-d)}$,
$\forall f\, Q_{1,2}^{f\prime}$, $Q_{5,6}^{(c-u,b,0)\prime}$,	$Q_6^{(0)}$	$Q_{5,6}^{(s-d)\prime}$, $Q_6^{(8d)\prime}$,
$Q_{1,2}^{(c-u,8d,0)}$, $Q_5^{(0)}$	$Q_{1,2}^{(b)}$	$Q_{5,6}^{(s-d,c-u,8d,b)}$,

3 2HDM

As a first example of an explicit NP model that can account for the measurement of ΔA_{CP}, we consider a two-Higgs-doublet model (2HDM), where a second scalar doublet,

$$\Phi \sim (1,2)_{-1/2} = \begin{pmatrix} \phi^0 \\ \phi^- \end{pmatrix}, \tag{20}$$

is added to the SM. A contribution to ΔA_{CP} arises if ϕ^0 couples to $u\bar{u}$ and $c\bar{u}$, generating both $D^0 \to K^+K^-$ and $D^0 \to \pi^+\pi^-$. Since all couplings besides $u\bar{u}$ and $c\bar{u}$ are irrelevant to this analysis, we take a conservative approach, considering minimal examples where Φ couples to u_R and is aligned with a single down-type LH mass eigenstate. This allows us to evade tree-level scalar mediated FCNC in the down sector. Assuming alignment with the quark doublet that has b_L as its down-type quark, we have [12]

$$\mathcal{L}_\Phi = -V(\Phi) + 2\lambda \left[\phi^0 \overline{U_{Li}} V_{ib} u_R + \phi^- \overline{b_L} u_R + \text{h.c.}\right], \tag{21}$$

where $U_{L1,2,3} = u_L, c_L, t_L$. Thus, the neutral scalar ϕ^0 couples u_R to u_L and c_L: $\lambda V_{cb}\phi^0 \bar{c}_L u_R + \lambda V_{ub}\phi^0_u \bar{u}_L u_R$. Integrating out the ϕ^0 field, these couplings lead to the effective four-quark coupling.

$$-\frac{8|\lambda|^2}{m^2_{\phi^0}} V_{ub}V^*_{cb}(\bar{u}_R c_L)(\bar{u}_L u_R) = \frac{|\lambda|^2}{m^2_{\phi^0}} V_{ub}V^*_{cb} Q^u_6. \tag{22}$$

The contribution to ΔA_{CP}, using Eq. (10), can be written as

$$\begin{aligned} \Delta A^\phi_{CP} &\approx \frac{2\sqrt{2}}{4|V_{us}V_{cs}|}\frac{G_0}{G_F}\mathcal{I}m(V_{ub}V^*_{cb})\mathcal{I}m(\Delta R^\phi) \\ &= \frac{\sqrt{2}}{4}\frac{G_0}{G_F} I_{\text{CKM}} \times \mathcal{I}m(\Delta R^\phi) \end{aligned} \tag{23}$$

where

$$G_0 \equiv 4|\lambda|^2/m^2_{\phi^0}, \tag{24}$$

and I_{CKM} is defined in Eq. (7). What is needed then to account for (3) is

$$\frac{G_0}{G_F} \simeq \frac{3.3}{\mathcal{I}m(\Delta R^\phi)} \implies \mathcal{I}m(\Delta R^\phi) G_0 \simeq \frac{1}{(160 \text{ GeV})^2}. \tag{25}$$

Thus, for $\mathcal{I}m(\Delta R^\phi) \in \{0.2 - 2\}$, we need $G_0^{-1/2} = m_{\phi^0}/(2|\lambda|) \in \{70, 230\}$ GeV.

3.1 Constraints from $D^0 - \overline{D}^0$ mixing

The scalar exchange contributes to $D^0 - \overline{D}^0$ mixing via box diagrams. Requiring that this contribution is not larger than the experimental constraints from Δm_D gives [12]

$$\frac{|\lambda|^4}{32\pi^2}\left(\frac{100 \text{ GeV}}{m_{\phi^0_u}}\right)^2 (V_{ub}V^*_{cb})^2 < 7 \times 10^{-9}, \tag{26}$$

or, equivalently

$$\frac{|\lambda|^2 G_0}{G_F} < 3 \times 10^3, \tag{27}$$

so, taking into account (25), the new contribution is negligible, allowing for the required G_0/G_F to explain ΔA_{CP}.

3.2 Constraints from ϵ'/ϵ

The same Yukawa couplings of ϕ^0 that contribute to direct CP violation in D decays, contribute unavoidably also to direct CP violation in K decays. The former effect comes at tree level and modifies ΔA_{CP}. The latter effect comes via box diagrams, involving ϕ^0 and a W-boson, and modifies ϵ'/ϵ. Upon integration out of ϕ^0 and W, we obtain the following effective four-quark coupling:

$$\frac{\sqrt{2}|\lambda|^2 G_F}{\pi^2}\Big(f_{uu}(x_\phi) - 2f_{ut}(x_\phi) + f_{tt}(x_\phi)\Big)V_{td}^* V_{ts}|V_{tb}|^2(\bar{d}_L u_R)(\bar{u}_R s_L), \tag{28}$$

where $x_\phi \equiv m_{\phi^0}^2/m_W^2$, and the loop function is given by

$$f_{ij}(x_\phi) = \frac{x_i^2 \log x_i}{(1-x_i)(x_j-x_i)(x_\phi-x_i)} + \frac{x_j^2 \log x_j}{(1-x_j)(x_i-x_j)(x_\phi-x_j)} + \frac{x_\phi^2 \log x_\phi}{(1-x_\phi)(x_i-x_\phi)(x_j-x_\phi)}. \tag{29}$$

Using the relation $(\bar{d}_L u_R)(\bar{u}_R s_L) = -\frac{1}{8}(\bar{d}_\alpha s_\beta)_{V-A}(\bar{u}_\beta u_\alpha)_{V+A} = -\frac{1}{8}Q_6^{u(ds)}$, we read off the corresponding Wilson coefficient,

$$C_6^{u(ds)} = -\frac{|\lambda|^2}{4\pi^2}\Big(f_{uu}(x_\phi) - 2f_{ut}(x_\phi) + f_{tt}(x_\phi)\Big)V_{td}^* V_{ts}|V_{tb}|^2. \tag{30}$$

Following Ref. [16], we use

$$\mathcal{R}e\left(\frac{\epsilon'}{\epsilon}\right) = -\frac{\omega}{\sqrt{2}|\epsilon|}\left(\frac{\mathcal{I}m(A_0)}{\mathcal{R}e(A_0)} - \frac{\mathcal{I}m(A_2)}{\mathcal{R}e(A_2)}\right), \tag{31}$$

and

$$\frac{\mathcal{I}m A_2^\phi}{\mathcal{R}e A_2} \approx \frac{3}{2}\frac{m_K^2}{m_s^2(m_c) - m_d^2(m_c)}\frac{\mathcal{I}m[\Delta C_6(m_c) + \frac{1}{3}\Delta C_5(m_c)]B_8^{(2)}(m_c)}{0.363|V_{us}^* V_{ud}|}, \tag{32}$$

where $\Delta C_i = C_i^u - C_i^d$. At the matching scale, our model generates $\Delta C_6(m_{\phi^0}) = C_6^u(m_{\phi^0})$, and $\Delta C_5(m_{\phi^0}) = 0$. Taking the conservative bound $\mathcal{R}e(\epsilon'/\epsilon)^\phi < \mathcal{R}e(\epsilon'/\epsilon)^{\rm Exp} \approx 1.66 \times 10^{-3}$, we reach the constraint

$$`C_6^{u(ds)}(m_{\phi^0}) < 2.23 \times 10^{-7}. \tag{33}$$

Figure 1 presents the various constraints together with curves for which Eq. (25) is satisfied with three representative values taken for $\mathcal{I}m(\Delta R^\phi)$. We conclude that ΔA_{CP} can be explained within this model, depending on the value of $\mathcal{I}m(\Delta R^\phi)$. For $\mathcal{I}m(\Delta R^\phi) \approx 1$, the mass of the neutral scalar is bounded to be $m_\phi \lesssim 235$ GeV, while for $\mathcal{I}m(\Delta R^\phi) \approx 0.2$ it is bounded to be very light, and subject to further constraints. For $\mathcal{I}m(\Delta R^\phi) \gtrsim 1.5$, the mass is unconstrained.

We note the following points:

- Two alternative choices for the Yukawa matrices such that only one down-type mass eigenstate is involved exist, with Φ aligned with the doublet containing either d_L or s_L. These suffer from large contributions to $D^0 - \overline{D}^0$ mixing, and therefore cannot account for ΔA_{CP}.
- It may seem surprising that this model can account ΔA_{CP} even though it contributes via the operator Q_6^{c-u}, disfavored by the EFT analysis. This is explained by the existence of additional contributions within this model to ϵ/ϵ', which interfere destructively. These are not taken into account in the EFT approach. Therefore this model evades the EFT conclusions regardless of the mass scale of the new scalars.
- We note that mid-range masses for the charged scalar (450 GeV $\lesssim m_{\phi^-} \lesssim$ a few TeV) are constrained by LHC dijet searches [17–19]. These would result in a further constraint in the $(|\lambda|, m_{\phi^0})$ plane, depending on the mass splitting between the neutral and charged scalars. Charged scalar masses below 450 GeV or above a few TeV are not constrained by these bounds.

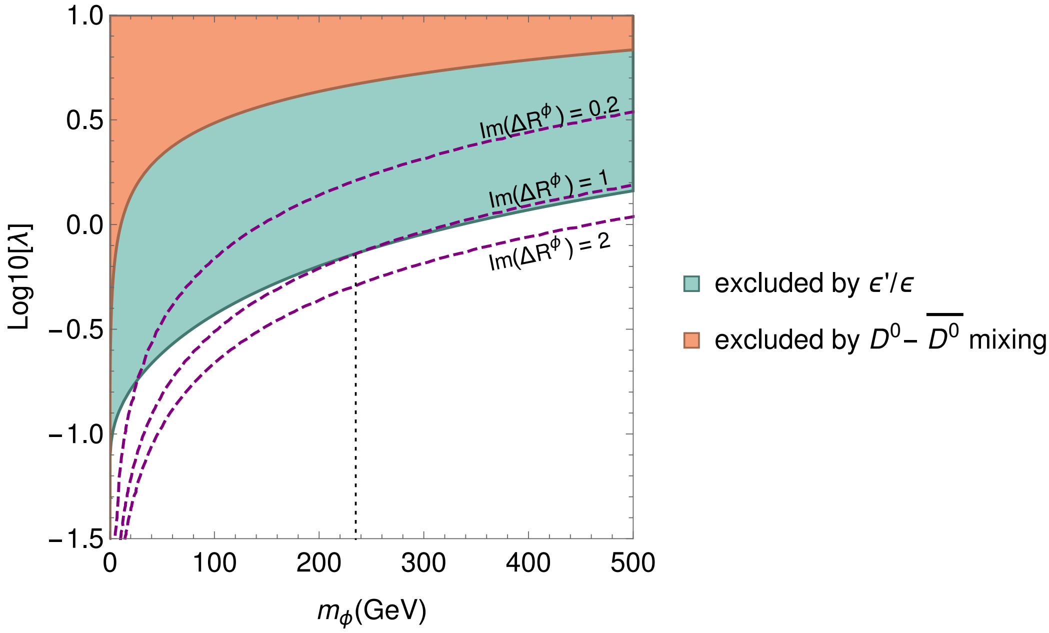

Figure 1: Excluded regions in parameter space due to $D^0 - \overline{D}^0$ mixing and ϵ'/ϵ constraints. The dashed lines depict the curves for which ΔA_{CP} is explained for $\mathcal{I}m(\Delta R^{\phi}) = 0.2, 1, 2$. The dotted vertical line marks the intersection, at $m_\phi \simeq 235\,\text{GeV}$.

4 MSSM

As a second example for candidate NP models to explain the measurement of ΔA_{CP}, we consider the MSSM. The dominant supersymmetric contribution to ΔA_{CP} is likely to come from loops involving gluinos and up-squarks. These contribute to the chromomagnetic operators Q_8 and Q'_8, which are very weakly constrained by $D^0 - \overline{D}^0$ mixing and ϵ/ϵ'. The dominant source of CP violation is likely to be the chirality-changing and flavor-changing mass-squared insertion [13],

$$\delta_{LR} \equiv (\delta^u_{LR})_{12} = \frac{(\tilde{M}^{2u}_{LR})_{12}}{\tilde{m}^2}, \tag{34}$$

where $\tilde{m}^2$ is the average up-squark mass, and $\tilde{M}^{2u}_{LR}$ is the left-right block in the 6×6 up-squark mass-squared matrix. In the approximation that only two squark generations are involved, we can express this parameter in terms of the supersymmetric mixing angles, $(K^u_{L,R})_{ij}$ and the mass-squared splitting between the squarks, $\Delta\tilde{m}^2_{ij}$:

$$\delta_{LR} = \frac{\Delta\tilde{m}^2_{q_{L1}q_{R2}}}{\tilde{m}^2}(K^u_L)_{12}(K^u_R)_{22}. \tag{35}$$

One can estimate the supersymmetric contribution as [13]

$$\Delta A_{CP} = 1.5 \times 10^{-3}\frac{\mathcal{I}m(\delta_{LR})}{2.5 \times 10^{-4}}\frac{1\text{ TeV}}{\tilde{m}} \times \mathcal{I}m(\Delta R^{\text{SUSY}}). \tag{36}$$

Thus in order to explain ΔA_{CP} we require

$$\mathcal{I}m(\delta_{LR}) \approx 2.5 \times 10^{-4}\frac{\tilde{m}}{1\text{ TeV}}\mathcal{I}m(\Delta R^{\text{SUSY}})^{-1}. \tag{37}$$

In MFV models [14],

$$\delta_{LR} \propto \frac{m_c}{\tilde{m}}(y_s^2 V_{us}V_{cs}^* + y_b^2 V_{ub}V_{cb}^*) \lesssim 10^{-7}, \tag{38}$$

and the contribution is negligible. In Froggatt-Nielsen (FN) models [14,20],

$$\delta_{LR} \sim \frac{\tilde{a}}{\tilde{m}}\frac{m_c|V_{us}|}{\tilde{m}} \sim 3 \times 10^{-4}\frac{\tilde{a}}{\tilde{m}}\frac{1\ \text{TeV}}{\tilde{m}}, \tag{39}$$

where $\tilde{a}$ is the typical scale of the trilinear scalar coupling. When comparing Eq. (39) to Eq. (37), it seems that FN-SUSY models are plausible candidates to account for ΔA_{CP}. One has to take into account, however, the FN relations with other entries of the squark mass-squared matrices, and, in particular,

$$\frac{\mathcal{I}m(\delta^u_{LR})_{12}}{\mathcal{I}m(\delta^q_{LR})_{11}} \sim \frac{m_c|V_{us}|}{m_q}, \quad (q = u, d). \tag{40}$$

Assuming phases of order one (which we do to explain ΔA_{CP}), the flavor-diagonal parameters are bounded by the EDM constraints. The resulting bounds are [14]

$$\begin{aligned} (\delta^u_{LR})_{12} &\lesssim 3 \times 10^{-4}\frac{\tilde{m}}{\text{TeV}} \quad (\text{from } (\delta^u_{LR})_{11}), \\ (\delta^u_{LR})_{12} &\lesssim 8 \times 10^{-5}\frac{\tilde{m}}{\text{TeV}} \quad (\text{from } (\delta^d_{LR})_{11}). \end{aligned} \tag{41}$$

Comparing to Eq. (36), we see that within FN, $\mathcal{I}m(\Delta R^{\text{SUSY}}) \gtrsim 3$ is required in order to explain ΔA_{CP}. In more elaborate flavor schemes (as in, for example, Ref. [21]) it is possible that Eq. (37) is satisfied for $\mathcal{I}m(\Delta R^{\text{SUSY}}) \approx 2$.

5 Vector-like quarks

A third example for a model that may explain ΔA_{CP} is a model exhibiting flavor changing Z couplings. Models with extra non-sequential quarks generally induce such flavor changing couplings for the Z boson. For example, the addition of vector-like up quarks in the $(3, 1, +2/3) \bigoplus (\bar{3}, 1, -2/3)$ representation induces flavor changing Z couplings of the form [2]

$$-\mathcal{L}_Z = \frac{gU^u_{ij}}{2\cos\theta_W}\overline{u_{L}}_i\gamma_\mu u_{Lj}Z^\mu + \text{h.c.} \tag{42}$$

The relevant coupling for ΔA_{CP} is U^u_{cu}, which also contributes at tree level to Δm_D, and at loop level to ϵ'/ϵ.

5.1 Constraint from $D^0 - \overline{D}^0$ mixing

The constraint from Δm_D can be calculated using the effective operators of Ref. [8]. The relevant $\Delta c = 2$ operator is $(\bar{u}_L\gamma^\mu c_L)^2 = \frac{1}{4}Q_1^{cu}$. Using Eq. (16) for the current bound on $\text{Re}(C_1^{cu})$, we arrive at

$$|U^u_{cu}| \lesssim 2.8 \times 10^{-4}. \tag{43}$$

5.2 Constraints from ϵ'/ϵ

A contribution to ϵ'/ϵ arises via a W-loop, inducing the operators $Q_{1,5}^{u(ds)} = (\bar{u}u)_{V\mp A}(\bar{s}d)_{V-A}$. We calculate the relevant Wilson coefficients and arrive at

$$\begin{aligned} C_1^{u(ds)} &= \frac{(3-4s_W^2)U_{cu}}{96\pi^2} \approx U_{cu}\cdot 2.2\times 10^{-3}, \\ C_5^{u(ds)} &= \frac{s_W^2 U_{cu} V_{cs} V_{ud}^*}{24\pi^2} \approx U_{cu}\cdot 9.3\times 10^{-4}. \end{aligned} \tag{44}$$

Using Eq. (18), the constraint on these coefficients is given by

$$\begin{aligned} \mathcal{I}m(C_1^{(c-u)(ds)}) &\lesssim 9.8\times 10^{-6}; \\ \mathcal{I}m(C_5^{(c-u)(ds)}) &\lesssim 7.1\times 10^{-7}, \end{aligned} \tag{45}$$

The constraint on $C_5^{(c-u)(ds)}$ is more stringent, implying

$$\mathcal{I}m(U_{cu}) \lesssim 7.6\times 10^{-4}. \tag{46}$$

ΔA_{CP} arises in this model through the tree level annihilation diagram $\bar{c}u \to \bar{u}u$, which contributes to the $\Delta c = 1$ four quark operators,

$$\begin{aligned} (\bar{u}_L\gamma_\mu c_L)(\bar{u}_L\gamma^\mu u_L) &= \frac{1}{4}Q_1^u, \\ (\bar{u}_L\gamma_\mu c_L)(\bar{u}_R\gamma^\mu u_R) &= \frac{1}{4}Q_5^u. \end{aligned} \tag{47}$$

The coefficients of these operators in this model are given by

$$\begin{aligned} C_1^u &= U_{cu}^u\left(\frac{1}{2}-\frac{2}{3}\sin^2\theta_W\right), \\ C_5^u &= U_{cu}^u\frac{2}{3}\sin^2\theta_W. \end{aligned} \tag{48}$$

Using Eq. (10), the contribution to ΔA_{CP} can be written as

$$\Delta A_{CP} \approx \frac{2}{|V_{us}V_{cs}|}\left(\mathcal{I}m(C_1^u)\mathcal{I}m(\Delta R_1^Z) + \mathcal{I}m(C_5^u)\mathcal{I}m(\Delta R_5^Z)\right). \tag{49}$$

when we have taken $\mathcal{I}m(\Delta R_1^Z) \approx \mathcal{I}m(\Delta R_5^Z) \equiv \mathcal{I}m(\Delta R^Z)$. Thus in order to explain the measurement we require

$$\mathcal{I}m(U_{cu}^u) \approx 1.84\times 10^{-4}\left(\frac{2}{\mathcal{I}m(\Delta R^Z)}\right), \tag{50}$$

which (under the assumption of $\mathcal{I}m(\Delta R^Z)\approx 2$) is allowed by Eqs. (43,46).

We note that this model is viable despite the fact that it induces the EFT-disfavored operator $Q_5^{(c-u)}$ (see Table 3), as its contibution to ΔA_{CP} is subleading to that of the operator $Q_1^{(c-u)}$.

6 Discussion

We have addressed the question of how easily can the new measurement of ΔA_{CP} be explained using benchmark NP models. We have followed the assumption that no significant hadronic enhancements are present, and derived the constraints coming mainly from measurements of $D^0 - \overline{D}^0$ mixing and ϵ'/ϵ. We find that non-generic though still simple NP models can account for the measured asymmetry.

Three candidate NP models were discussed – 2HDM, MSSM and vector-like up-quarks. Our assumption of no significant hadronic enhancements is implemented by allowing at most $\mathcal{I}m(\Delta R^{\rm SM,NP}) \approx 2$, in our Eq. (10). We find that:

- Both a 2HDM where scalar $(c\bar{u}),(u\bar{u})$ couplings are present and models with vector-like up-quarks inducing $(c\bar{u})$ Z couplings can account for the measured asymmetry.
- The MSSM combined with flavor frameworks (MFV, FN) is unable to produce the desired contribution (FN requires $\mathcal{I}m(\Delta R^{\rm FN}) \gtrsim 3$). The MSSM with a generic flavor structure is unconstrained.

Ref. [6] studied the scenario where the SM accounts for ΔA_{CP} with mild $SU(3)$ breaking effects but a strong enhancement of $\Delta U = 0$ transitions. They obtain two predictions: U-spin invariant strong phases should be large, and $A_{CP}(K^+K^-) \approx -A_{CP}(\pi^+\pi^-)$. Interestingly, in all three models that we analyzed the new physics operators that account for ΔA_{CP} do not introduce new sources of U-spin breaking, and thus the latter prediction does not favor the SM over these models.

In all three specific new physics models, the flavor structure is not in the minimal flavor violation class, and in fact it is non-generic. Thus, it is difficult to make definite predictions for the modification of other flavor changing and/or CP violating processes. Yet, it is unlikely that the only significant modification would be to singly Cabibbo suppressed charm decays. This situation motivates a broad flavor precision program, such as in the LHCb and BELLE-II experiments.

Of course, a direct search for the new degrees of freedom required by the various models is also well motivated. The upper bound on the scale of new physics is model dependent, and varies from few tens of TeV in the low energy EFT, to hundreds of GeV in the 2HDM.

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