PROCEEDINGS OF SCIENCE

Charm decays

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I discuss hadronic decays of D mesons with emphasis on the recent discovery of charm CP violation in $D^0 \to K^+K^-, \pi^+\pi^-$ decays. The measured difference $\Delta a_{CP} \equiv a_{CP}^{\rm dir}(D^0 \to K^+K^-) - a_{CP}^{\rm dir}(D^0 \to \pi^+\pi^-) = (-15.4 \pm 2.9) \cdot 10^{-4}$ of two direct CP asymmetries exceeds the SM prediction by a factor of 7. A possible explanation is an enhancement of the penguin amplitude entering $a_{CP}^{\rm dir}$ by QCD effects which are not understood yet. Alternatively, Δa_{CP} could be dominated by contributions from new physics. In order to distinguish these two hypotheses further CP asymmetries should be measured. To this end CP asymmetries resulting from the interference of two tree-level amplitudes auch as $a_{CP}^{\rm dir}(D^0 \to K_S K_S)$ or $a_{CP}^{\rm dir}(D^0 \to \overset{(-)}{K}{}^{*0} K_S)$ are especially interesting.

18th International Conference on B-Physics at Frontier Machines - Beauty2019 -
29 September – 4 October, 2019
Ljubljana, Slovenia

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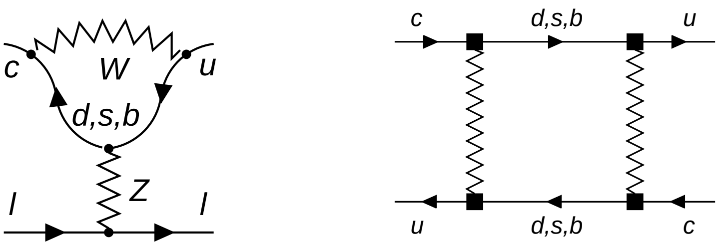


Figure 1: FCNC amplitudes: Z penguin diagram contributing to $D^0 \to \ell^+\ell^-$ (left) and D–$\bar{D}$ mixing box diagram (right).

1. Overview

The charm event of the year 2019 was the announcement of March 21,

LHCb sees a new flavour of matter-antimatter asymmetry,

presenting the first observation of CP violation in charm decays. The LHCb collaboration has measured the difference of two direct CP asymmetries [1]:

$$\begin{aligned} \Delta a_{CP} \equiv\ & a_{CP}^{\rm dir}(D^0 \to K^+K^-) - a_{CP}^{\rm dir}(D^0 \to \pi^+\pi^-) \\ =\ & (-15.4 \pm 2.9) \cdot 10^{-4}. \end{aligned} \tag{1.1}$$

Before discussing the theory aspects of this measurement I give a short overview on the role of charm decays in particle physics and the methods and difficulties of theory predictions. While weak decays of charmed hadrons are not useful for the metrology of the Cabibbo-Kobayashi-Maskawa (CKM) matrix, they have a unique role in probing new physics in the flavour sector of up-type quarks. Flavour-changing neutral current (FCNC) amplitudes (see Fig. 1 for examples) involve the CKM combinations $\lambda_d = V_{cd}^* V_{ud}$, $\lambda_s = V_{cs}^* V_{us}$, and $\lambda_b = V_{cb}^* V_{ub}$ associated with d, s, and b quarks, respectively, on internal lines of the FCNC loop diagrams. CKM unitarity $\lambda_d + \lambda_s + \lambda_b = 0$ allows us to eliminate one of these CKM combinations. If we write $p \equiv \sum_q \lambda_q p(m_q)$ for the penguin diagram in Fig. 1 and choose to eliminate λ_d, we find $p = \lambda_s[p(m_s) - p(m_d)] + \lambda_b[p(m_b) - p(m_d)]$. The loop contribution with λ_b is tiny because of $|\lambda_b| \sim 10^{-4}$, while the contribution proportional to $\lambda_s \simeq \lambda = 0.22$ vanishes in the limit $m_d = m_s$ (corresponding to unbroken U-spin symmetry) and is therefore heavily suppressed by the Glashow-Iliopoulos-Maiani (GIM) mechanism. The latter feature also makes it impossible to predict FCNC processes in a reliable way. For example, a perturbative calculation of the loop function $p(m_s) - p(m_d)$ involving internal d and s quarks in the penguin diagram of Fig. 1 gives a result proportional to

$$\frac{G_F}{M_Z^2} \cdot \underbrace{\underbrace{(m_s - m_d)}}_{\substack{\text{U-spin}\\ \text{breaking}\\ \text{GIM}}} \cdot \underbrace{\underbrace{(m_s + m_d)}}_{\substack{\text{artefact of}\\ \text{perturbation theory}}} \tag{1.2}$$

The presence of small quark masses below the QCD scale $\Lambda_{\rm QCD} \sim 400\,\text{MeV}$ indicates that the perturbative calculation is not trustworthy. While the factor $m_s - m_d$ correctly catches the linear

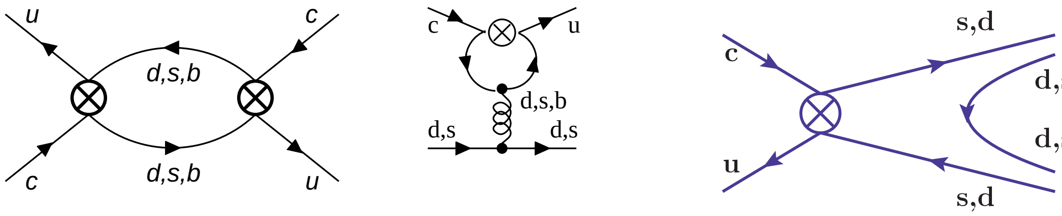


Figure 2: Box diagram describing D–$\bar{D}$ mixing, penguin diagram contributing to SCS decays, and tree-level exchange diagrams contributing to CP asymmetries in decays into two neutral kaons. The cross denotes a W propagator contracted to a point as in the Fermi theory.

U-spin breaking of the amplitude, the factor $m_s + m_d$ occurs, because the left-chiral nature of the W coupling requires an even number of left-right flips on the internal quark line. This factor is an artefact of perturbation theory, non-perturbative QCD provides other sources of left-right flips, for instance the quark condensate. The only experimentally established FCNC transition in charm physics is the D–$\bar{D}$ mixing amplitude, the mass and width difference between the two neutral D eigenstates (normalized to the total width Γ) are (HFLAV) [2]

$$x = \frac{\Delta m}{\Gamma} = 0.39^{+0.11}_{-0.12}\%, \qquad y = \frac{\Delta\Gamma}{2\Gamma} = 0.651^{+0.063}_{-0.069}\%. \tag{1.3}$$

These numbers exceed the naive perturbative result of the box diagram in Fig. 1 by far. Still our theoretical understanding of D–$\bar{D}$ mixing is too poor to conclude whether the measurements in Eq. (1.3) involve new physics contributions or not. Thus while charm FCNC transitions are highly suppressed in the Standard Model (SM), our insufficient understanding of low-energy QCD effects limits their use as new-physics analyzers.

The other avenue to new physics are measurements of CP asymmetries. Hadronic weak decays of the D mesons

$$D^+ \sim c\bar{d}, \qquad D^0 \sim c\bar{u}, \qquad D_s^+ \sim c\bar{s}, \tag{1.4}$$

are denoted Cabibbo-favored (CF) or singly or doubly Cabibbo-suppressed (SCS or DCS), if the decay amplitude is proportional to λ^0, λ^1, or λ^2, respectively. Non-zero CP asymmetries require the interference of two amplitudes with different CP-violating phases, which implies that all SM predictions for charm CP asymmetries involve the suppression factor $\mathrm{Im}\,\frac{\lambda_b}{\lambda_s} = -6 \cdot 10^{-4}$. We may categorize the detectable CP asymmetries by their origin from

- *box–tree*,
- *penguin–tree*, or
- *tree–tree*

interference. The first category contains mixing-induced CP asymmetries like $a_{CP}^{\mathrm{mix}}(D^0(t) \to K^+\pi^-)$. Most direct CP asymmetries are in the second category and the LHCb measurement in Eq. (1.1) has established penguin-tree interference, if interpreted within the SM. The CP asymmetries of the third category arise from the interference of the $c \to us\bar{s}$ and $c \to ud\bar{d}$ amplitudes. Fig. 2 shows sample diagrams for the three categories of CP asymmetries.

Theoretical studies of weak decays of charmed hadrons heavily utilize the approximate $\mathrm{SU(3)_F}$ symmetry of QCD. The QCD lagrangian is invariant under unitary rotations of the light quark triplet (u,d,s) in the limit $m_u = m_d = m_s$. The SU(2) subgroup of unitary rotations of (u,d) is the strong isospin (I-spin) symmetry; the counterpart for (s,d) is the above-mentioned U-spin symmetry and V-spin refers to rotations of (s,u). The I-spin breaking of QCD scales like $(m_d - m_u)/\Lambda_{\mathrm{QCD}} \sim 0.02$, while U-spin holds to to an accuracy of order $(m_s - m_d)/\Lambda_{\mathrm{QCD}} \sim 0.3$.

2. CP violation in penguin-tree interference

It is helpful to use $\lambda_d + \lambda_s + \lambda_b = 0$ to decompose the amplitude $\mathscr{A}^{\mathrm{SCS}}$ of the SCS decay of a charged or neutral D meson into two light mesons M, M' as [3]

$$\mathscr{A}^{\mathrm{SCS}}(MM') \equiv \lambda_{sd} A_{sd}(MM') - \frac{\lambda_b}{2} A_b(MM') \tag{2.1}$$

with

$$\lambda_{sd} = \frac{\lambda_s - \lambda_d}{2} \qquad \text{and} \qquad -\frac{\lambda_b}{2} = \frac{\lambda_s + \lambda_d}{2}. \tag{2.2}$$

If we write the effective hamiltonian as

$$H \;=\; \lambda_d H_d + \lambda_s H_s + \lambda_b H_b + \mathrm{h.c.} \tag{2.3}$$

with

$$H_q \;=\; 4\frac{G_F}{\sqrt{2}} \left[C_1\, \bar{u}_L^\alpha \gamma_\mu c_L^\alpha\, \bar{q}_L^\beta \gamma^\mu q_L^\beta \;+\; C_2\, \bar{u}_L^\alpha \gamma_\mu c_L^\beta\, \bar{q}_L^\beta \gamma^\mu q_L^\alpha \right], \tag{2.4}$$

where G_F is the Fermi constant, then

$$A_{sd}(MM') = \langle MM'|H_s - H_d|D\rangle, \qquad A_b(MM') = \langle MM'|H_s + H_d - 2H_b|D\rangle. \tag{2.5}$$

H_q contains the Wilson coefficients $C_{1,2}$ with the perturbative QCD corrections to the W exchange diagram. $C_{1,2}$ multiply the four-quark operators describing the W-mediated weak interaction (where the Fierz relation $\bar{q}_L \gamma_\mu c_L \bar{u}_L \gamma^\mu q_L = \bar{u}_L \gamma_\mu c_L \bar{q}_L \gamma^\mu q_L$ is used) and α, β are color indices.

The benefit of the decomposition in Eq. (2.1) becomes clear from Eq. (2.5): A_{sd} is a $|\Delta U| = 1$ amplitude, because $H_s - H_d$ involves $\bar{s}s - \bar{d}d$ which transforms like a U-spin triplet. Likewise A_b is a $\Delta U = 0$ amplitude. In view of $|\lambda_b|/|\lambda_{sd}| \sim 10^{-3}$ we can work to first non-vanishing order in λ_b and may safely replace $\lambda_{sd} = \lambda_s + \lambda_b/2$ by λ_s. Data on branching ratios can be used to determine $|A_{sd}|$ for the decay modes of interest, but are not accurate enough to give information on $|A_b|$.

To discuss the direct CP asymmetry in some SCS decay $D \to MM'$ we need Eq. (2.1) for $\mathscr{A} = \mathscr{A}^{\mathrm{SCS}}(MM')$ and the analogous decomposition for the amplitude $\overline{\mathscr{A}}$ of the CP-conjugate decay $\bar{D} \to \bar{M}\bar{M}'$, where $\mathrm{CP}\,|D\rangle = -|\bar{D}\rangle$ and $\mathrm{CP}\,|MM'\rangle = |\bar{M}\bar{M}'\rangle$:

$$\overline{\mathscr{A}} = -\lambda_{sd}^* A_{sd} + \frac{\lambda_b^*}{2} A_b.$$

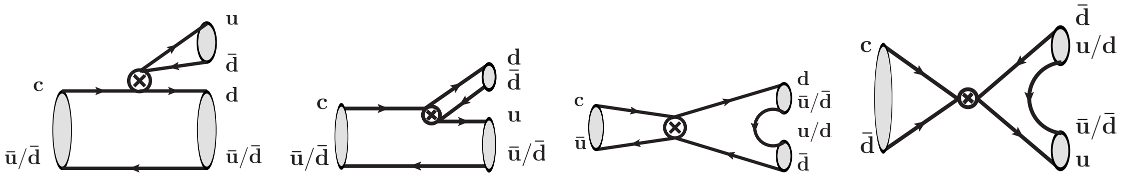


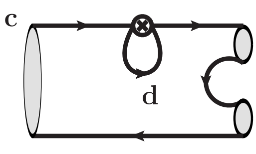
Figure 3: Topological amplitudes *color-favored tree* T, *color-suppressed tree* C, *exchange* E, and *annihilation* A. They are understood to include all perturbative and non-perturbative strong-interaction effects, i.e. one may view the diagrams as dressed with an arbitrary number of gluons. T, C, E, and A are complex numbers which are determined from a global fit to data. .

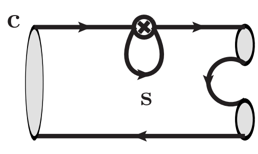
The SM prediction for the desired CP asymmetry reads

$$a_{CP}^{\rm dir} \equiv \frac{|\mathscr{A}|^2 - |\overline{\mathscr{A}}|^2}{|\mathscr{A}|^2 + |\overline{\mathscr{A}}|^2} = \operatorname{Im}\frac{\lambda_b}{\lambda_{sd}}\operatorname{Im}\frac{A_b}{A_{sd}} = -6\cdot 10^{-4}\operatorname{Im}\frac{A_b}{A_{sd}}. \qquad (2.6)$$

One can conveniently describe $\bar{D} \to \bar{M}\bar{M}'$ decays in terms of topological amplitudes [4, 5]. In the SU(3)$_{\rm F}$ limit we can express A_{sd} of all $D^0, D^+, D_s^+ \to MM'$ decays for all combinations $M, M' = \pi^{\pm,0}, K^{\pm,0}$ as linear combinations of the four tree diagrams T, C, E, and A shown in Fig. 3. Linear (i.e. first-order) SU(3)$_{\rm F}$ breaking can be included in the method in a straightforward way [6] . The topological-amplitude method is mathematically equivalent [7] to the decomposition of the decay amplitudes in terms of matrix elements classified by their SU(3)$_{\rm F}$ symmetry properties [8, 9]. A global fit of all branching ratios to the four SU(3)$_{\rm F}$ limit amplitudes of Fig. 3 returns a poor fit. If one includes the topological amplitudes parametrising linear SU(3)$_{\rm F}$ breaking the fit is underconstrained and one obtains a perfect fit on a large submanifold of the parameter space [7]. By assuming upper bounds on the sizes of the SU(3)$_{\rm F}$-breaking topological amplitudes (limiting their magnitudes to e.g. 30% of the leading T amplitude) one can nevertheless derive useful constraints on T, C, E, A and the SU(3)$_{\rm F}$-breaking amplitudes [7].

The information from branching ratios is not sufficient to predict CP asymmetries: A_b in Eq. (2.6) involves new topological amplitudes in the SU(3)$_{\rm F}$ limit, which cannot be constrained from branching fractions. These are the *penguin* amplitude P and the *penguin annihilation* amplitude PA. Consider

$P_d \equiv$ 

$P_s \equiv$ 

and the analogously defined P_b. The amplitude A_b of a SCS decay involves

$$P \equiv P_d + P_s - 2P_b \qquad (2.7)$$

and/or the analogous combination $PA \equiv PA_d + PA_s - 2PA_b$ defined in terms of the PA amplitude in Fig. 4. P and PA are $\Delta U = 0$ amplitudes and therefore do not appear in A_{sd} constrained from branching ratio data.

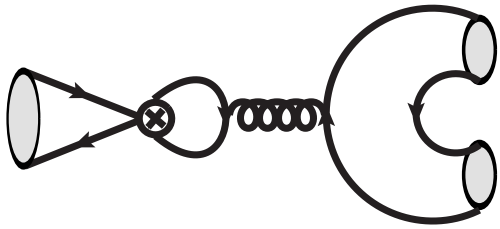


Figure 4: Topological *penguin annihilation* amplitude PA_q, where q is the quark flavor of the loop. The perturbative gluon provides the hard momentum transfer from the loop to the final state, further soft QCD interaction is needed to arrange the correct color quantum numbers.

In the SU(3)$_F$ limit one finds $A_b(\pi^+\pi^-) = A_b(K^+K^-)$, $A_{sd}(\pi^+\pi^-) = -A_{sd}(K^+K^-)$ [10], and

$$\text{Im}\,\frac{A_b(\pi^+\pi^-)}{A_{sd}(\pi^+\pi^-)} = -\text{Im}\,\frac{A_b(K^+K^-)}{A_{sd}(K^+K^-)} = \text{Im}\,\frac{P+PA}{A_{sd}(\pi^+\pi^-)}. \tag{2.8}$$

In the last equation $A_b(K^+K^-) = A_{sd}(K^+K^-) + P + PA$ [11] has been used.

A consequence of the SU(3)$_F$ relation in Eq. (2.8) for Δa_{CP} in Eq. (1.1) is

$$\Delta a_{CP} \overset{\text{SU(3) limit}}{=} 2a_{CP}^{\text{dir}}(D^0 \to K^+K^-). \tag{2.9}$$

Thus in the SM we expect Δa_{CP} to be twice as large as the individual CP asymmetries, up to corrections from SU(3)$_F$ breaking. $a_{CP}^{\text{dir}}(D^0 \to K^+K^-) = -a_{CP}^{\text{dir}}(D^0 \to \pi^+\pi^-)$ is an example of an SU(3)$_F$ sum rule relating different CP asymmetries to each other [8]. One can improve such sum rules by including first-order breaking SU(3)$_F$ breaking in A_{sd} and e.g. find a refined sum rule involving the direct CP asymmetries in $D^0 \to K^+K^-$, $D^0 \to \pi^+\pi^-$, and $D^0 \to \pi^0\pi^0$ [7]. This is possible, because the global fit on D branching ratios returns information on magnitudes and phases of the topological amplitudes contributing to A_{sd} for the three amplitudes.

The history of measurements of Δa_{CP} prior to the 2019 discovery is as follows:

Previous LHCb measurements:
2011 [12]: $\Delta a_{CP} = (-82 \pm 21 \pm 11) \cdot 10^{-4}$
2014 [13]: $\Delta a_{CP} = (-14 \pm 16 \pm 8) \cdot 10^{-4}$
2016 [14]: $\Delta a_{CP} = (-10 \pm 8 \pm 3) \cdot 10^{-4}$

Previous world averages (HFLAV):
2015: $\Delta a_{CP} = (-25.3 \pm 10.4) \cdot 10^{-4}$
2016: $\Delta a_{CP} = (-13.4 \pm 7.0) \cdot 10^{-4}$

Theoretical analyses of CP asymmetries based on SU(3)$_F$ symmetry can relate different CP asymmetries but cannot predict the overall size because of the a priori unknown P and PA amplitudes. In 2011 LHCb presented the first evidence for a non-zero Δa_{CP} with the value quoted above [12], which was unexpectedly large. All SU(3)$_F$ papers written afterwards (such as Ref. [11]) present ranges for Δa_{CP} compatible with the value in Eq. (1.1), because they use the 2011 value as input. This feature merely reflects the fact that Δa_{CP} in Eq. (1.1) complies with the earlier measurement presented in Ref. [12].

Confronting

$$a_{CP}^{\text{dir}}(D^0 \to K^+K^-) \simeq \frac{1}{2}\Delta a_{CP} = \frac{1}{2}(-15.4 \pm 2.9) \cdot 10^{-4} \tag{2.10}$$

with Eq. (2.6) one can solve for the imaginary part of the "penguin-to-tree ratio":

$$\frac{1}{2}\frac{A_b(K^+K^-)}{A_{sd}(K^+K^-)} \approx \frac{P_d}{A_{sd}(K^+K^-)} \tag{2.11}$$

to find [15]

$$\frac{1}{2}\frac{A_b(K^+K^-)}{A_{sd}(K^+K^-)} = 0.65 \pm 0.12. \tag{2.12}$$

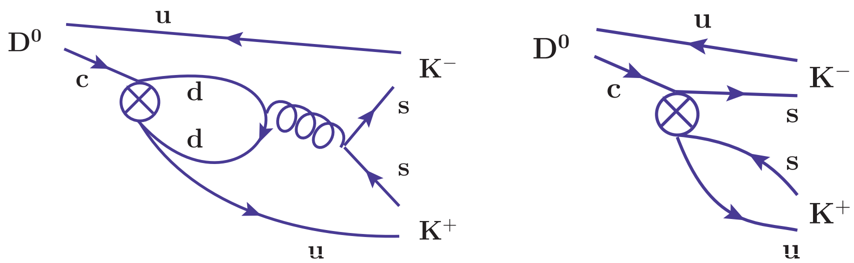
Methods employing a perturbative calculation of the penguin diagram in Fig. 1 give much smaller values for the ratio in Eq. (2.11). The authors of Ref. [15] conclude that there is either a non-perturbative enhancement mechanism of the $\Delta U = 0$ amplitude A_b (i.e. an enhancement of $P + PA$) [3] or physics beyond the SM (BSM).

The momentum flowing through the penguin loop in P and PA are of order 1 GeVor larger, therefore a perturbative calculation of this loop is not unreasonable. In QCD sum rule calculations this loop is indeed calculated perturbatively and Ref. [16] finds

$$|\Delta a_{CP}| \leq (2.0 \pm 0.3) \cdot 10^{-4}, \tag{2.13}$$

which is smaller than the experimental value by a factor of 7! QCD sum rules are a well established method successfully describing many quantities in B physics while poorly tested in D physics. An essential ingredient of QCD sum rule calculations is the assumptions that the combined effect of all highly excited hadronic resonances and multi-hadron states is correctly described by a perturbative calculation.

Next I argue that one arrives at an estimate in the ballpark of Eq. (2.13) even without invoking a perturbative treatment of the penguin loop. We need $\mathrm{Im}\frac{A_b}{A_{sd}} = \frac{\mathrm{Im}A_b A_{sd}^*}{|A_{sd}|^2}$ and the numerator $\mathrm{Im}A_b A_{sd}^*$ is the absorptive part of the penguin-tree interference term:



By the optical theorem this absorptive part is related to a $c \to ud\bar{d}$ decay followed by $d\bar{d} \to s\bar{s}$ rescattering. This rescattering is essential for a non-zero direct CP asymmetry and we may discuss it without referring to the perturbative picture of quarks and gluons. One contribution is $D^0 \to \pi^+\pi^- \to K^+K^-$ rescattering. Each such contribution to $\mathrm{Im}\frac{A_b}{A_{sd}}$ is color-suppressed $\propto 1/N_c$ and further suppressed by a factor of $\sim 1/\pi$ from the phase space integral of the rescattering process. We conclude that we need an enhancement factor X for the $\Delta U = 0$ transitions feeding A_b such that $X \cdot \frac{1}{N_c \pi} \stackrel{!}{=} 0.65 \pm 0.12$. This means $X \sim 6$, thus the QCD sum rule result of Ref. [16] has the expected size and is not unnaturally small. A resonant enhancement involving only the $\Delta U = 0$ channel leaves A_{sd} unchanged and can therefore accomodate Δa_{CP} in Eq. (1.1) without violating data on branching fractions which comply with the SM [7]. In Ref. [17] it has been suggested that the $f_0(1710)$ resonance (having a mass close to the D^0 mass) could provide such an enhancement mechanism through $D^0 \to f_0(1710) \to K^+K^-$ or $\pi^+\pi^-$. For this mechanism to work the overlap of the $f_0(1710)$ state with the K^+K^- or $\pi^+\pi^-$ state must be sufficiently large, in contradiction with the expectation that a high resonance will dominantly decay to high-multiplicity states. More insight will be gained from measurements of the branching fractions of $f_0(1710)$ into K^+K^- or $\pi^+\pi^-$. Since in $D^0 \to f_0(1710) \to MM'$ decays the final state carries the quantum numbers of the $f_0(1710)$ one can find $SU(3)_F$ relations among different CP asymmetries which are specific to this mechanism and may serve to falsify the $f_0(1710)$ resonance hypothesis [17].

Physics beyond the SM may well affect Δa_{CP}. If the BSM contribution to $c \to ud\bar{d}$ or $c \to us\bar{s}$ comes with an arbitrary $\mathscr{O}(1)$ CP phase, the suppression factor $\mathrm{Im}\,\frac{\lambda_b}{\lambda_{sd}} = -6 \cdot 10^{-4}$ is absent and the exchange of a virtual multi-TeV particle could induce a Δa_{CP} in the range of Eq. (1.1). Various BSM scenarios with heavy particles are discussed in Ref. [18]. Also light BSM particles with feeble couplings may explain the measured Δa_{CP}; Ref. [19] studies a model with a Z' boson. If the new physics couples differently to s and d quarks (i.e. if it violates U-spin symmetry), then $a_{CP}^{\rm dir}(K^+K^-) \approx -a_{CP}^{\rm dir}(\pi^+\pi^-)$ does not hold. Thus such new-physics scenarios can be distinguished from the hypothesis of QCD enhanced A_b amplitudes. To this end one must measure one of the individual CP asymmetries $a_{CP}^{\rm dir}(K^+K^-)$ and $a_{CP}^{\rm dir}(\pi^+\pi^-)$ or their sum.

3. CP violation in tree-tree interference

Whenever the tree-level transitions $c \to u\bar{d}d$ and $c \to u\bar{s}s$ interfere, the decay can have a non-vanishing direct CP asymmetry proportional to

$$\mathrm{Im}\,\frac{V_{ud}V_{cd}^*}{V_{us}V_{cs}^*} = \mathrm{Im}\,\frac{-V_{us}V_{cs}^* - V_{ub}V_{cb}^*}{V_{us}V_{cs}^*} = -\mathrm{Im}\,\frac{V_{ub}V_{cb}^*}{V_{us}V_{cs}^*} \simeq -\mathrm{Im}\,\frac{\lambda_b}{\lambda_{sd}} \simeq 6 \cdot 10^{-4}. \tag{3.1}$$

Tree-tree interference occurs for final states containing an $\eta^{(\prime)}, \omega, \ldots$, or a pair of neutral Kaons like $K_SK_S, K_SK^{*0}, \ldots$, or for multibody final states like $K^+K^-\pi^+\pi^-$ containing all four $s, \bar{s}, d, \bar{d}$ quarks. The topological amplitudes E (in Fig. 2 on the right) and PA (in Fig. 4) constitute A_b entering the CP asymmetry in D^0 decays into two neutral Kaons. The global fit to two-body D^0, D^+, D_s^+ decays into two pseudoscalars in Ref. [7] has returned a large value of E, so that $A_b(K_SK_S)$ and $a_{CP}^{\rm dir}(K_SK_S)$ in the SM can be large [20]:

$$|a_{CP}^{\rm dir}(D^0 \to K_SK_S)| \leq 1.1\% \qquad @95\%\ \text{C.L.} \tag{3.2}$$

Throughout this talk it is assumed that the CP violation in Kaon mixing is properly subtracted from the measured $a_{CP}^{\rm dir}$ [22]. The ratio $A_b(K_SK_S)/A_{sd}(K_SK_S)$ is large, because $A_{sd}(K_SK_S)$ vanishes in the $SU(3)_F$ limit, while $A_b(K_SK_S)$ does not. The size of the $D^0 \to K_SK_S$ branching fraction (proportional to $|A_{sd}|^2$) measures the size of $SU(3)_F$ breaking in E [7, 20]. The maximal value in Eq. (3.2) corresponds to the maximal value of $|2E+PA|$ returned by the fit of Ref. [7] in addition to a favorable strong phase difference $\arg(A_b/A_{sd}) = \pm\pi/2$. More likely values for $|a_{CP}^{\rm dir}(D^0 \to K_SK_S)|$ are three times smaller than the upper bound in Eq. (3.2). If the strong phase $\arg(A_b/A_{sd})$ is close to zero, $|a_{CP}^{\rm dir}(D^0 \to K_SK_S)|$ will be too small to be measured. However, in this case one will find instead a larger mixing-induced CP asymmetry in $D^0(t) \to K_SK_S$ [20].

Other interesting decay modes to study CP violation from tree–tree interference are $D^0 \to \bar{K}^{*0}K_S$ and $D^0 \to K^{*0}K_S$. Since the final state is not a CP eigenstate, these decay modes offer more possibilities for CP studies. As a special feature of these modes the CP asymmetry persists even in the untagged sample of $\overset{(-)}{D} \to \bar{K}^{*0} K_S$ and one can determine a non-vanishing CP asymmetry by just counting $\overset{(-)}{D} \to \bar{K}^{*0} K_S$ and $\overset{(-)}{D} \to K^{*0} K_S$ events [21] in a sample with equal number of D^0 and $\bar{D}^0$ decays. In real life, however, one must study the four Dalitz plots of $D^0, \bar{D}^0 \to (K^-\pi^+)_{\bar{K}^{*0}} K_S$ and $D^0, \bar{D}^0 \to (K^+\pi^-)_{K^{*0}} K_S$ to take care of interferences with other decay modes leading to a $K^\mp \pi^\pm K_S$ final state.

The SM prediction is [21]

$$|a_{CP}^{\rm dir}(D^0 \to \bar{K}^{*0} K_S)| \leq 0.003, \tag{3.3}$$

and the same bound applies to $|a_{CP}^{\rm dir}(D^0 \to K^{*0} K_S)|$. In the SU(3)$_{\rm F}$ limit $a_{CP}^{\rm dir}(D^0 \to \bar{K}^{*0} K_S) = -a_{CP}^{\rm dir}(D^0 \to K^{*0} K_S)$ holds. The value in Eq. (3.3) is smaller than the one in Eq. (3.2), because $A_{sd}(\bar{K}^{*0} K_S)$ and $A_{sd}(K^{*0} K_S)$ do not vanish in the SU(3)$_{\rm F}$ limit. The prediction in Eq. (3.3) uses data from an LHCb analysis of the $D^0 \to K^\mp \pi^\pm K_S$ Dalitz plot [23].

The original motivation to study CP violation in tree-tree interference was the possibility of large CP asymmetries in the SM, i.e. the $D^0 \to K_S K_S$ and $D^0 \to \overset{(-)}{K}{}^{*0} K_S$ modes were proposed as discovery channels for CP violation in charm decays [20, 21]. Now, in view of the experimental result in Eq. (1.1) the measurement of CP asymmetries from tree-tree interference will instead give valuable insight into the mechanism underlying the large value in Eq. (1.1). For example, QCD dynamics enhancing P and PA by a factor of 7 cannot enhance $|a_{CP}^{\rm dir}(D^0 \to K_S K_S)|$ or $|a_{CP}^{\rm dir}(D^0 \to K_S K_S)|$ and $|a_{CP}^{\rm dir}(D^0 \to \overset{(-)}{K}{}^{*0} K_S)|$ over the results in Eqs. (3.2) and (3.3) by the same factor of 7. In Sec. V of Ref. [20] the correlation of the imprints of new physics on various CP asymmetries is discussed.

4. Summary

All CP asymmetries in the SM are proportional to the small factor ${\rm Im}\, \frac{\lambda_b}{\lambda_{sd}} \simeq -6 \cdot 10^{-4}$, which makes these asymmetries sensitive to new physics. The measured value in Eq. (1.1) exceeds the theory prediction [16] by a factor of 7. An explanation within the SM calls for enhanced QCD effects in $\Delta U = 0$ transitions [3, 15] whose origin is currently not understood. With more precise data on other charm CP asymmetries we can hope to find out whether a QCD effect or BSM physics is behind Δa_{CP} in Eq. (1.1) [10, 11, 18–21]. This discrimination will be straightforward, if the new physics couples differently to d and s quarks, so that the SU(3)$_{\rm F}$ sum rules of Refs. [8, 11] are violated beyond the expected level of SU(3)$_{\rm F}$ breaking.

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