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Theory of D-meson decays

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I review the theory of D-meson decays, with a focus on mixing and CP violation. Strategies to identify new-physics contributions to direct CP violation are also discussed.

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1. Introduction

Charm physics is an important and interesting part of flavor physics and has caused some excitement recently. The dominance of the first two quark generations and an efficient Glashow-Iliopoulos-Maiani (GIM) mechanism leads to generically small standard-model effects for mixing and CP violation. However, the large long-distance contributions are difficult to evaluate. Thus D-meson decays are less suited for precision extractions of Cabibbo-Kobayashi-Maskawa (CKM) matrix elements than, for instance, K- and B-meson decays. On the other hand, the neutral decays are interesting because they form a bound state of up-type quarks which exhibits particle-antiparticle mixing, and new-physics effects could easily be of the same order of magnitude as the small standard-model contribution.

For instance, it was widely assumed that the size of direct CP violation in D decays is small within the standard model. By the end of 2011, the LHCb collboration published a somewhat unexpectedly large value for the difference of CP asymmetries in the $D^0 \to K^+K^-$ and $D^0 \to \pi^+\pi^-$ modes, $\Delta\mathscr{A}_{CP} \equiv \mathscr{A}_{CP}(D \to K^+K^-) - \mathscr{A}_{CP}(D \to \pi^+\pi^-)$, namely $\Delta\mathscr{A}_{CP} = (-0.82 \pm 0.21 \pm 0.11)\%$ [1]. This measurement was then soon confirmed by the CDF collaboration, who found $\Delta\mathscr{A}_{CP} = (-0.62 \pm 0.21 \pm 0.10)\%$ [2]. Subsequently, some theoretical groups argued that this value could be accommodated within the standard model. At the beginning of this year, the LHCb collaboration updated their original pion-tagged analysis to find the significantly smaller value $\Delta\mathscr{A}_{CP} = (-0.34 \pm 0.15 \pm 0.10)\%$ [3], and published a second, muon-tagged analysis, yielding $\Delta\mathscr{A}_{CP} = (+0.49 \pm 0.30 \pm 0.14)\%$ (note the unexpected positive sign) [4]. Inclusion of BaBar [5] and Belle [6] results leads to the current world average $\Delta\mathscr{A}_{CP} = (-0.329 \pm 0.121)\%$ provided by the HFAG group. While this new value can likely be accommodated within the standard model, the large theoretical uncertainties do not allow us to exclude a new-physics contribution.

In this article we will discuss the theory of hadronic two-body D decays, with an emphasis on mixing and CP violation, and look at $\Delta\mathscr{A}_{CP}$ in more detail. We will also mention some strategies to look for effects of new physics in direct CP violation.

2. Hadronic two-body D-meson decays

Flavor-changing transitions are induced in the standard model by the exchange of W bosons, whose couplings to quarks are proportional to CKM matrix elements. Because of the hierarchy of the CKM matrix, D-meson decays are dominated by the first two generations of quarks, up to corrections of order λ^4, where $\lambda \equiv |V_{us}| \approx 0.23$. We thus expect CP-violating effects to be small. We can classify the two-body decay modes according to the size of the contributing CKM matrix elements as Cabibbo-favored (CF), singly (SCS) and doubly (DCS) Cabibbo-suppressed. Examples are the modes $D^0 \to K^-\pi^+$ (proportional to λ^0); $D^0 \to \pi^+\pi^-$, $D^0 \to K^+K^-$, $D^0 \to K^0\bar{K}^0$, $D^+ \to \pi^+\pi^0$ (suppressed by one power of λ); and $D^0 \to K^+\pi^-$ (suppressed by two powers of λ).

The starting point for most analyses is the weak effective Hamiltonian [7]. The effective field theory formalism allows to separate short- and long-distance contributions. The former, parameterized by the Wilson coefficients of the effective operators, can be calculated in perturbation theory with, in principle, arbitrary precision, including the summation of large-logarithmic contributions to all orders in the strong coupling constant. The hadronic matrix elements, on the other hand, are dominated by nonperturbative QCD effects and much harder to evaluate. The heavy-quark expansion is not expected to work well, due to the large value of $\Lambda_{\rm QCD}/m_c$. Flavor symmetries ($SU(3)_{\rm flavor}$ and its subgroups like isospin and U spin), plus breaking effects, can help to extract hadronic matrix elements from data, with possibly large uncertainties. The ultimate way will be the computation of the matrix elements using lattice QCD, probably in the not so near future [8].

3. Mixing and CP violation

The neutral D-meson system is unique in that it exhibits particle-antiparticle mixing in the up-quark sector: top quarks decay before they can hadronize, and up quarks do not form neutral bound states with opposite quantum numbers.

The short-distance contribution to D mixing is heavily CKM-suppressed, and the remaining contributions of internal light quarks represent non-local effects of exchanged light mesons, which cannot be evaluated perturbatively. In general, the time evolution in the Wigner-Weisskopf approximation of the neutral D-meson system is given by the Schrödinger equation

$$i\frac{d}{dt}\begin{pmatrix}|D(t)\rangle\\|\bar{D}(t)\rangle\end{pmatrix}=\left(M-i\frac{\Gamma}{2}\right)\begin{pmatrix}|D(t)\rangle\\|\bar{D}(t)\rangle\end{pmatrix}, \tag{3.1}$$

where M and Γ are the Hermitean mass- and decay matrices, respectively. The equation is solved by going to the mass eigenstate basis

$$|D_{H,L}\rangle = p|D^0\rangle \mp q|\bar{D}^0\rangle \tag{3.2}$$

which diagonalizes the Hamiltonian (see [9] for details), with eigenvalues

$$M_{H,L} - i\Gamma_{H,L}\,. \tag{3.3}$$

They are conventionally combined into the parameters

$$\Gamma_D \equiv \frac{\Gamma_H+\Gamma_L}{2}\,,\quad x\equiv\frac{M_H-M_L}{\Gamma_D}\,,\quad y\equiv\frac{\Gamma_H-\Gamma_L}{2\Gamma_D}\,. \tag{3.4}$$

The subscripts H and L stand for the heavier and lighter mass eigenstate, respectively. The theory prediction for x and y is difficult. They vanish for exact $SU(3)$ flavor symmetry [10]; non-zero value arise only at second order in flavor-symmetry breaking [11]. We can express these parameters as sums over on-shell or off-shell intermediate states:

$$\begin{aligned} y &= \frac{1}{2\Gamma_D}\sum_n \rho_n[\langle D^0|\mathscr{H}|n\rangle\langle n|\mathscr{H}|\bar{D}^0\rangle + \langle\bar{D}^0|\mathscr{H}|n\rangle\langle n|\mathscr{H}|D^0\rangle]\,,\\ x &= \frac{1}{\Gamma_D}\left[\langle D^0|\mathscr{H}|\bar{D}^0\rangle + P\sum_n\frac{\langle D^0|\mathscr{H}|n\rangle\langle n|\mathscr{H}|\bar{D}^0\rangle + \langle\bar{D}^0|\mathscr{H}|n\rangle\langle n|\mathscr{H}|D^0\rangle}{M_D^2-E_n^2}\right]. \end{aligned} \tag{3.5}$$

Here ρ_n is the density of the state n and P denotes the principle value. The first (local) term in the second line is negligible in the standard model; however, it could get contributions from heavy new particles above the electroweak scale. There are two main approaches to predict x and y within the standard model [11], with both their advantages and shortcomings.

The inclusive approach [10, 12, 13] amounts to expanding Eq. (3.5) in a series of local operators of increasing dimension. The leading dimension-six box diagrams yield the estimate $x_{\rm box} \sim$ few $\times 10^{-5}$, $y_{\rm box} \sim$ few $\times 10^{-7}$, where $x_{\rm box} \propto m_s^4$, $y_{\rm box} \propto m_s^6$. These chiral suppressions can be lifted by quark condensates in higher-order contributions. Including dimension-eight operators, the estimates are $x \sim y \lesssim 10^{-3}$ [11]. Note that using heavy-quark expansion, the authors of [14] do not exclude values of $y \sim 1\%$ in the standard model. All these values constitute order-of-magnitude estimates rather than presise predictions. The assumed so-called quark-hadron duality holds exactly only in the limit $m_c \to \infty$ and might receive large corrections.

In the exclusive approach [15, 11, 16] y is evaluated by performing the sum over exclusive final states, taking into account, among others, effects of $SU(3)$ breaking and resonances. While the authors of [16] find $x \sim y \lesssim 10^{-3}$, in [11] values of $y \sim 1\%$ appear naturally. y can then be related to x via a dispersion relation, yielding $x \sim 1\%$ [17].

We conclude that there are large uncertainties in the standard-model prediction of the mixing parameters in the neutral D-meson system. Thus, usually the experimental values, which are both of order 0.5%, are used as constraints on new physics, demanding that new contributions at most saturate the measured values within the experimental uncertainties.

Let us now take into account CP violation. As is well known, we can distinguish three types of CP violation in meson decays (for details, see the review [9]): CP violation in decay, CP violation in mixing, and CP violation in the interference of decays with and without mixing. Looking at the time-integrated CP asymmetry for final the CP eigenstate f

$$a_f \equiv \frac{\Gamma(D^0 \to f) - \Gamma(\bar{D}^0 \to f)}{\Gamma(D^0 \to f) + \Gamma(\bar{D}^0 \to f)} = a_f^d + a_f^m + a_f^i\,, \tag{3.6}$$

we see that it contains all three types of CP violation, here denoted by a_f^d, a_f^m, a_f^i, respectively. Using the experimental constraints on x, y, we find the simplified expressions [18]:

$$a^m = -\frac{y}{2}\left(\left|\frac{q}{p}\right| - \left|\frac{p}{q}\right|\right)\cos\phi_D\,, \quad a^i = \frac{x}{2}\left(\left|\frac{q}{p}\right| + \left|\frac{p}{q}\right|\right)\sin\phi_D\,, \tag{3.7}$$

where ϕ_D is the phase of $\lambda_f \equiv \frac{q}{p}\frac{\bar{A}_f}{A_f}$, and $A_f = A(D^0 \to f)$ is the amplitude for a D meson decaying into a final state f. Indirect CP violation (a_f^m, a_f^i), related to mixing effects, is expected to be very small in the standard model and is universal (process-independent) to a good approximation.

Let us now focus on direct CP violation. We can write the relevant amplitudes as

$$A_f = A_f^T\left[1 + r_f e^{i(\delta_f - \phi_f)}\right], \quad \bar{A}_f = A_f^T\left[1 + r_f e^{i(\delta_f + \phi_f)}\right], \tag{3.8}$$

where r_f is the relative magnitude of the subleading penguin amplitude with a relative strong phase δ_f and a weak phase ϕ_f. Both of them have to be nonzero in order to have nonvanishing direct CP violation. This can be seen in the expression

$$a_f^d := \frac{|A_f|^2 - |\bar{A}_f|^2}{|A_f|^2 + |\bar{A}_f|^2} = 2 r_f \sin\phi_f \sin\delta_f\,, \tag{3.9}$$

where we kept only the leading term in the small ratio r_f. The contribution $a^m + a^i$ of indirect CP violation is universal and cancels to good approximation in

$$\Delta\mathscr{A}_{CP} := a^d_{K^+K^-} - a^d_{\pi^+\pi^-}\,. \tag{3.10}$$

The world average for $\Delta\mathscr{A}_{CP}$ quoted in the introduction still constitutes an almost 3σ evidence for non-zero direct CP violation in $D \to K^+K^-, \pi^+\pi^-$. The question is whether this is a clear signal for new physics.

An estimation of the hadronic matrix elements in the limit $m_c \to \infty$ yields a value of $\Delta\mathscr{A}_{CP}$ roughly a factor of three below the measurement. However, from $SU(3)$ fits [19, 20, 21, 22] we know that power corrections are large. This signals the breakdown of the $1/m_c$ expansion. The fact that penguin contraction matrix elements can be large has already been observed in [23]; they can account for the measured value of $\Delta\mathscr{A}_{CP}$ and decay rate difference, with nominal U-spin breaking of order 20% [24]. Various groups found at least marginal consistency of the old, larger value with the standard model [22, 25, 26, 27]. The new, smaller value can likely be accommodated within the standard model.

4. Testing for new physics

We have seen in the last section that $\Delta\mathscr{A}_{CP}$ can likely be accommodated within the standard model. Nevertheless, a new-physics contribution to $\Delta\mathscr{A}_{CP}$ is by no means excluded. In this section we describe some strategies to disentangle the two contributions.

If new physics resides in new contributions to $\Delta I = 1/2$ transitions (I stands for isospin), they can be disentangled using isospin sum rules. The basic observation is the following [28]: The standard-model tree-level effective Hamiltonian for $D \to \pi\pi$ has both $\Delta I = 1/2$ and $\Delta I = 3/2$ contributions, $Q_T \sim (\bar{d}c)(\bar{u}d)$. The QCD penguin operators are have only a pure $\Delta I = 1/2$ component, $Q_P \sim (\bar{c}u) \otimes (\bar{u}u + \bar{d}d + \bar{s}s)$. It follows that $\Delta I = 3/2$ direct CP-violating transitions are absent in SM, since they would need two interfering amplitudes.

There are isospin-breaking effects of $\mathscr{O}(1\%)$, about the same order of magnitude as the value of $\Delta\mathscr{A}_{CP}$. However, isospin breaking by quark masses and QED effects is CP conserving. The CP-violating contribution of the electroweak penguins is down by $\alpha/\alpha_s \approx 1\%$ with respect to the already small QCD-penguin contribution.

The isospin decomposition of the amplitudes is

$$A_{\pi^+\pi^-} = \frac{1}{\sqrt{6}}\mathscr{A}_{3/2} + \frac{1}{\sqrt{3}}\mathscr{A}_{1/2}\,, \quad A_{\pi^0\pi^0} = \frac{1}{\sqrt{3}}\mathscr{A}_{3/2} - \frac{1}{\sqrt{6}}\mathscr{A}_{1/2}\,, \quad A_{\pi^+\pi^0} = \frac{\sqrt{3}}{2}\mathscr{A}_{3/2}\,. \tag{4.1}$$

For instance, note that the decay $D^+ \to \pi^+\pi^0$ is pure $\Delta I = 3/2$ transition, so any CP asymmetry in this mode would be due to new physics. However, the converse is not true: The strong phase could be smaller between the standard-model and new-physics $\Delta I = 3/2$ amplitudes than between the $\Delta I = 1/2$ and $\Delta I = 3/2$ amplitudes. Another test is given by the following sum rule:

$$|A_{\pi^+\pi^-}|^2 - |\bar{A}_{\pi^-\pi^+}|^2 + |A_{\pi^0\pi^0}|^2 - |\bar{A}_{\pi^0\pi^0}|^2 - \frac{2}{3}(|A_{\pi^+\pi^0}|^2 - |\bar{A}_{\pi^-\pi^0}|^2) = \frac{1}{2}(|\mathscr{A}_{1/2}|^2 - |\bar{\mathscr{A}}_{1/2}|^2)\,. \tag{4.2}$$

If this sum is zero, while the individual rate differences are nonzero, CP asymmetries are likely dominated by new physics effects in the $\Delta I = 3/2$ transition.

Analogous rules apply for each polarization state of $D^+ \to \rho^+\rho^0$. It is possible to write down sum rules also for the $D \to \rho\pi$, $D \to K^{(*)}\bar{K}^{(*)}\pi(\rho)$, $D_s^+ \to K^*\pi(\rho)$ decay modes [28].

The isospin sum rules test only for new physics in $\Delta I = 3/2$ transitions. However, most well-motivated models give new contributions to the gluon penguin $Q_{8g} = \frac{m_c}{4\pi^2}\bar{u}_L\sigma_{\mu\nu}T^a g_s G_a^{\mu\nu} c_R$. A prominent example is left-right squark mixing in supersymmetric models [18, 29, 30]. δ_{LR} contributions to Q_{8g} are enhanced by $m_{\tilde{g}}/m_c$, whereas there is no such enhancement of $\Delta C = 2$ operators [18]. In this way we can have $\mathscr{O}(10^{-2})$ effects in direct CP violation and still evade bounds from $D^0 - \bar{D}^0$ mixing. For a more complete overview, see the review [31].

The measurement of radiative decay modes can help testing for $\Delta I = 1/2$ new physics [32]. The first key observation here is that new-physics models with a large chromo-magnetic penguin Q_{8g} typically also have large electromagnetic dipole operator $Q_{7\gamma} = \frac{m_c}{4\pi^2}\bar{u}_L\sigma_{\mu\nu}Q_u e F^{\mu\nu} c_R$. Even if this is not the case, a sizeable Wilson coefficient will be generated by QCD effects through operator mixing.

The second key observation is then that $Q_{7\gamma}$ has a large matrix element in the $D \to (V = \rho^0, \omega)\gamma$ and $D \to (\phi \to K^+K^-)\gamma$ decay modes. A detailed analysis [32] shows that we can expect asymmetries as large as

$$\begin{aligned} \approx 10\%\,, &\quad D \to \rho^0\gamma,\, D \to \omega\gamma; \\ \approx 2\%\,, &\quad D \to K^+K^-\gamma \quad \text{(above the } \phi \text{ resonance).} \end{aligned} \tag{4.3}$$

This is one order of magnitude above what can reasonably be expected within the standard model. The usual caveat is that strong phases might be small. This limitation could be overcome with a time-dependent analysis $D \to V\gamma$ or $D \to V(\gamma * \to \ell^+\ell^-)$, which is however experimentally challenging.

5. Conclusion

Charm physics is theoretically and experimentally challenging. Although D-meson decays are currently not suited for precision extraction of CKM-matrix elements, some observables are in principle very sensitive to new-physics contributions. In this article, we focused on a very small subset of interesting modes. There are many more interesting modes which have not been mentioned above. The decays $D^+ \to \ell^+\nu$ are helicity suppressed. For instance, $D^+ \to e^+\nu$ has a branching ratio of order $\mathscr{O}(10^{-8})$. The flavor-changing neutral current transitions $D^0 \to \mu^+\mu^-, e^+e^-, \gamma\gamma$ are long-distance dominated in the standard model. The corresponding branching ratios are estimated to be very small; a recent analysis finds $\mathrm{Br}(D^0 \to \mu^+\mu^-) \sim 3 \times 10^{-13}$, $\mathrm{Br}(D^0 \to \gamma\gamma) \sim 10^{-8}$ [33]. The branching ratios for the modes $D^0 \to \mu^\pm e^\mp$ are zero in the standard model to very good approximation. Any observation of these modes well above the standard-model expectations would be a clear signal of new physics.

Lastly we would like to mention that D-meson decays provide an important testbed for technical tools like lattice-QCD calculations or heavy-quark expansion. Furthermore, they can provide information also about the light-quark sector [34].

We expect that progress on the theoretical side, in particular from lattice QCD, as well as experimental data from LHCb and flavor factories, will give new input to our efforts to find new phenomena beyond the standard model using D-meson decays. Charm physics is full of surprises and will be an important part of the flavor-physics program also in the coming years.

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References

[1] R. Aaij *et al.* [LHCb Collaboration], Phys. Rev. Lett. **108**, 111602 (2012)

[2] T. Aaltonen *et al.* [CDF Collaboration], Phys. Rev. Lett. **109**, 111801 (2012)

[3] [LHCb Collaboration], LHCb-CONF-2013-003; CERN-LHCb-CONF-2013-003

[4] R. Aaij *et al.* [LHCb Collaboration], Phys. Lett. B **723**, 33 (2013)

[5] B. Aubert *et al.* [BaBar Collaboration], Phys. Rev. Lett. **100**, 061803 (2008)

[6] B. R. Ko [Belle Collaboration], arXiv:1212.1975 [hep-ex].

[7] G. Buchalla, A. J. Buras and M. E. Lautenbacher, Rev. Mod. Phys. **68**, 1125 (1996)

[8] M. T. Hansen and S. R. Sharpe, Phys. Rev. D **86**, 016007 (2012)

[9] K. Anikeev, D. Atwood, F. Azfar, S. Bailey, C. W. Bauer, W. Bell, G. Bodwin and E. Braaten *et al.*, hep-ph/0201071.

[10] H. Georgi, Phys. Lett. B **297** (1992) 353 [hep-ph/9209291].

[11] A. F. Falk, Y. Grossman, Z. Ligeti and A. A. Petrov, Phys. Rev. D **65**, 054034 (2002) [hep-ph/0110317].

[12] T. Ohl, G. Ricciardi and E. H. Simmons, Nucl. Phys. B **403** (1993) 605 [hep-ph/9301212].

[13] I. I. Y. Bigi and N. G. Uraltsev, Nucl. Phys. B **592** (2001) 92 [hep-ph/0005089].

[14] M. Bobrowski, A. Lenz, J. Riedl and J. Rohrwild, JHEP **1003** (2010) 009 [arXiv:1002.4794 [hep-ph]].

[15] A. AAnselm and Y. I. Azimov, Phys. Lett. B **85** (1979) 72; J. F. Donoghue, E. Golowich, B. R. Holstein and J. Trampetic, Phys. Rev. D **33** (1986) 179; P. Colangelo, G. Nardulli and N. Paver, Phys. Lett. B **242**, 71 (1990); T. A. Kaeding, Phys. Lett. B **357** (1995) 151 [hep-ph/9505393]; E. Golowich and A. A. Petrov, Phys. Lett. B **427** (1998) 172 [hep-ph/9802291].

[16] H. -Y. Cheng and C. -W. Chiang, Phys. Rev. D **81** (2010) 114020 [arXiv:1005.1106 [hep-ph]].

[17] A. F. Falk, Y. Grossman, Z. Ligeti, Y. Nir and A. A. Petrov, Phys. Rev. D **69**, 114021 (2004) [hep-ph/0402204].

[18] Y. Grossman, A. L. Kagan and Y. Nir, Phys. Rev. D **75** (2007) 036008 [hep-ph/0609178].

[19] H. -Y. Cheng and C. -W. Chiang, Phys. Rev. D **81** (2010) 074021 [arXiv:1001.0987 [hep-ph]].

[20] H. -Y. Cheng and C. -W. Chiang, Phys. Rev. D **85** (2012) 034036 [Erratum-ibid. D **85** (2012) 079903] [arXiv:1201.0785 [hep-ph]].

[21] B. Bhattacharya, M. Gronau and J. L. Rosner, Phys. Rev. D **85** (2012) 054014 [arXiv:1201.2351 [hep-ph]].

[22] D. Pirtskhalava and P. Uttayarat, Phys. Lett. B **712** (2012) 81 [arXiv:1112.5451 [hep-ph]].

[23] M. J. Savage, Phys. Lett. B **257**, 414 (1991).

[24] J. Brod, A. L. Kagan and J. Zupan, Phys. Rev. D **86** (2012) 014023 [arXiv:1111.5000 [hep-ph]].

[25] T. Feldmann, S. Nandi and A. Soni, JHEP **1206**, 007 (2012)

[26] E. Franco, S. Mishima and L. Silvestrini, JHEP **1205** (2012) 140 [arXiv:1203.3131 [hep-ph]].

[27] G. Hiller, M. Jung and S. Schacht, Phys. Rev. D **87** (2013) 014024 [arXiv:1211.3734 [hep-ph]].

[28] Y. Grossman, A. L. Kagan and J. Zupan, Phys. Rev. D **85** (2012) 114036 [arXiv:1204.3557 [hep-ph]].

[29] G. F. Giudice, G. Isidori and P. Paradisi, JHEP **1204** (2012) 060 [arXiv:1201.6204 [hep-ph]].

[30] G. Hiller, Y. Hochberg and Y. Nir, Phys. Rev. D **85** (2012) 116008 [arXiv:1204.1046 [hep-ph]].

[31] W. Altmannshofer, R. Primulando, C. -T. Yu and F. Yu, JHEP **1204** (2012) 049 [arXiv:1202.2866 [hep-ph]].

[32] G. Isidori and J. F. Kamenik, Phys. Rev. Lett. **109** (2012) 171801 [arXiv:1205.3164 [hep-ph]].

[33] A. Paul, I. I. Bigi and S. Recksiegel, Phys. Rev. D **82** (2010) 094006 [Erratum-ibid. D **83** (2011) 019901] [arXiv:1008.3141 [hep-ph]].

[34] G. Ricciardi, Phys. Rev. D **86** (2012) 117505 [arXiv:1209.3386 [hep-ph]].