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D^0–$\bar{D}^0$ mixing and new physics

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Abstract. The Standard Model (SM) and new physics (NP) descriptions of D^0 mixing are discussed. The SM part of the discussion addresses both quark-level and hadron-level contributions. The NP part describes our recent works on the rate difference $\Delta\Gamma_{\mathrm{D}}$ and the mass difference ΔM_{D}. In particular, we describe how the recent experimental determination of ΔM_{D} is found to place tightened restrictions on parameter spaces for 17 out of the 21 NP models considered.

Keywords. D meson; new physics; D^0–$\bar{D}^0$ mixing.

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1. Introduction

Every time the mixing in a neutral meson–antimeson system has been observed, we have learned something new and very important. For example, kaon mixing was first observed in 1956, and has opened the way to the study of the K_1–K_2 system, and the understanding of the general phenomenon of quantum mixing, eventually to neutrino oscillations on the one hand and the discovery of parity violation, and eventually CP violation, on the other. Finally, when theorists came up with the Standard Model, we understood the origin of the mixing as being due to the charm threshold and that the mass of the charmed quark had to be about 2 GeV, in 1974! The mixing in the B system was first observed in 1987, and found to be almost as large as in the K system, and it pointed to a rather heavy top quark, confirmed in 1995. B_s mixing was observed to be quite large in 2006, and confirmed our knowledge of the quark mixing matrix elements. The latest is the mixing in D system which has just been found at the order of a %, and it is possible that it too will lead to hints of new physics. That D mixing can be a probe of new physics was recognized very early, by Amitava Datta. He pointed out in 1985 [1] that D mixing in Standard Model (SM) will be very small, because in addition to the near perfect GIM cancellation, there is the additional suppression by a factor of $(m_s/m_c)^2$ due to the external quark being heavier than the internal lines in the box diagram; and he went on to discuss new physics scenarios [2] that could be probed with the observation of D mixing. It has taken us over 20 years to finally reach that stage.

The most recent values for D mixing are as summarized by heavy flavour averaging group (HFAG) [3],

$$x_{\rm D} \equiv \frac{\Delta M_{\rm D}}{\Gamma_{\rm D}} = \left(8.4^{+3.2}_{-3.4}\right) \cdot 10^{-3},$$
$$y_{\rm D} \equiv \frac{\Delta \Gamma_{\rm D}}{2\Gamma_{\rm D}} = (6.9 \pm 2.1) \cdot 10^{-3}. \tag{1}$$

The observed signal is seen to occur at about the 1% level. Allowing for the possibility of an NP component in D^0 mixing amplitude,

$$\mathcal{M}_{\rm mix} = \mathcal{M}_{\rm SM} + \mathcal{M}_{\rm NP}. \tag{2}$$

The relative phase between $\mathcal{M}_{\rm SM}$ and $\mathcal{M}_{\rm NP}$ is not known. Thus, in our detailed study of various NP contributions to $x_{\rm D}$ we usually compare the NP predictions to $\pm 1\sigma, \pm 2\sigma$ windows relative to the central $x_{\rm D}$ value of eq. (1).

1.1 *Operator product expansion (OPE) and renormalization group*

An important technical aspect of refs [4,5] is the process of relating an amplitude at some NP scale $\mu = M$ to one at, say, the charm scale $\mu = m_{\rm c}$. This takes the form

$$\langle f|\mathcal{H}_{\rm NP}|i\rangle = G \sum_{i=1} C_i(\mu) \langle f|\mathcal{Q}_i|i\rangle(\mu), \tag{3}$$

where the prefactor G has the dimension of inverse-squared mass, the C_i are dimensionless Wilson coefficients, and the $\mathcal{Q}_i$ are the effective operators. At the leading order of dimension six, it turns out that there are eight four-quark operators,

$$\begin{aligned}
\mathcal{Q}_1 &= (\bar{u}_{\rm L}\gamma_\mu c_{\rm L})(\bar{u}_{\rm L}\gamma^\mu c_{\rm L}) \\
\mathcal{Q}_2 &= (\bar{u}_{\rm L}\gamma_\mu c_{\rm L})(\bar{u}_{\rm R}\gamma^\mu c_{\rm R}), \\
\mathcal{Q}_3 &= (\bar{u}_{\rm L} c_{\rm R})(\bar{u}_{\rm R} c_{\rm L}), \\
\mathcal{Q}_4 &= (\bar{u}_{\rm R} c_{\rm L})(\bar{u}_{\rm R} c_{\rm L}), \\
\mathcal{Q}_5 &= (\bar{u}_{\rm R}\sigma_{\mu\nu} c_{\rm L})(\bar{u}_{\rm R}\sigma^{\mu\nu} c_{\rm L}), \\
\mathcal{Q}_6 &= (\bar{u}_{\rm R}\gamma_\mu c_{\rm R})(\bar{u}_{\rm R}\gamma^\mu c_{\rm R}), \\
\mathcal{Q}_7 &= (\bar{u}_{\rm L} c_{\rm R})(\bar{u}_{\rm L} c_{\rm R}), \\
\mathcal{Q}_8 &= (\bar{u}_{\rm L}\sigma_{\mu\nu} c_{\rm R})(\bar{u}_{\rm L}\sigma^{\mu\nu} c_{\rm R}).
\end{aligned} \tag{4}$$

Any given NP contribution will often involve several of these, but in any event never more than these eight. The evolution is determined by solving the RG equations obeyed by the Wilson coefficients,

$$\frac{\rm d}{{\rm d}\log\mu}\vec{C}(\mu) = \hat{\gamma}^T \vec{C}(\mu), \tag{5}$$

where $\hat{\gamma}$ is the 8×8 anomalous dimension matrix [6]. The output of this calculation is a set of RG factors $r_i(\mu, M)$ which are expressed in terms of ratios of QCD fine structure constants evaluated at different scales, e.g. as with

$$r_1(\mu, M) = \left(\frac{\alpha_s(M)}{\alpha_s(m_t)}\right)^{2/7} \left(\frac{\alpha_s(m_t)}{\alpha_s(m_b)}\right)^{6/23} \left(\frac{\alpha_s(m_b)}{\alpha_s(\mu)}\right)^{6/25}. \tag{6}$$

1.2 *Operator matrix elements*

One needs ultimately to evaluate the D^0–$\bar{D}^0$ matrix elements of the eight operators $\{\mathcal{Q}_i\}$. In general, eight non-perturbative parameters would need to be evaluated by some means such as a lattice determination. As a practical matter, the method used in refs [4,5] is to introduce a 'modified vacuum saturation' (MVS), where all such matrix elements are written in terms of the known matrix elements of $(V-A) \times (V-A)$ and $(S-P) \times (S+P)$ matrix elements $B_{\rm D}$ and $B_{\rm D}^{\rm (S)}$ [6],

$$\begin{aligned}
\langle \mathcal{Q}_1 \rangle &= \frac{2}{3} f_{\rm D}^2 M_{\rm D}^2 B_{\rm D}, \\
\langle \mathcal{Q}_2 \rangle &= -\frac{1}{2} f_{\rm D}^2 M_{\rm D}^2 B_{\rm D} - \frac{1}{N_{\rm c}} f_{\rm D}^2 M_{\rm D}^2 \bar{B}_{\rm D}^{\rm (S)}, \\
\langle \mathcal{Q}_3 \rangle &= \frac{1}{4N_{\rm c}} f_{\rm D}^2 M_{\rm D}^2 B_{\rm D} + \frac{1}{2} f_{\rm D}^2 M_{\rm D}^2 \bar{B}_{\rm D}^{\rm (S)}, \\
\langle \mathcal{Q}_4 \rangle &= -\frac{2N_{\rm c}-1}{4N_{\rm c}} f_{\rm D}^2 M_{\rm D}^2 \bar{B}_{\rm D}^{\rm (S)}, \\
\langle \mathcal{Q}_5 \rangle &= \frac{3}{N_{\rm c}} f_{\rm D}^2 M_{\rm D}^2 \bar{B}_{\rm D}^{\rm (S)}, \\
\langle \mathcal{Q}_6 \rangle &= \langle \mathcal{Q}_1 \rangle, \\
\langle \mathcal{Q}_7 \rangle &= \langle \mathcal{Q}_4 \rangle, \\
\langle \mathcal{Q}_8 \rangle &= \langle \mathcal{Q}_5 \rangle,
\end{aligned} \tag{7}$$

where the number of colours $N_{\rm c} = 3$ and, as in refs [7,8], we define

$$\bar{B}_{\rm D}^{(S)} \equiv B_{\rm D}^{\rm (S)} \cdot \frac{M_{\rm D}^2}{(m_c + m_u)^2}. \tag{8}$$

With the above theoretical machinery in hand, we are now ready to consider SM and NP contributions to D^0 mixing.

2. Standard Model analysis

One can use quarks or hadrons as the basic degrees of freedom in carrying out the SM analysis of D^0 mixing. In principle, these should give the same result. However, as we shall see, rather different features appear in each description.

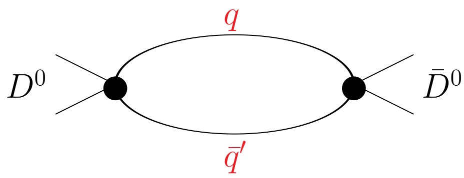

Figure 1. Loop diagram for $D^0 \to \bar{D}^0$.

Table 1. Flavour cancellations in $\Delta\Gamma_{\rm D}$.

Intermediate state	$\mathcal{O}(z^0)$	$\mathcal{O}(z^1)$	$\mathcal{O}(z^2)$
$s\bar{s}$	1/2	$-3z$	$3z^2$
$d\bar{d}$	1/2	0	0
$s\bar{d} + d\bar{s}$	-1	$3z$	$-3z^2$
Total	0	0	0

2.1 *Quark-level analysis*

At leading order in the SM, the OPE for D^0 mixing consists of two dimension-six four-quark operators [9]. The next order contains 15 dimension-nine six-quark operators. For each increasing order in the OPE, there are still more local quark and gluon operators and the problem of determining operator matrix elements becomes ever more severe. For this reason, the dimension-six sector has received by far the most attention.

Since the b-quark is essentially decoupled due to the tiny V_{ub} value, only the light d, s quarks propagate in the loop. In the limit of $m_d \to 0$, the only mass ratio that appears in the problem is

$$z \equiv (m_s/m_c)^2 \simeq 0.006. \tag{9}$$

Table 1 examines one of the loop-functions for $\Delta\Gamma_{\rm D}$ and shows the results of carrying out an expansion in powers of z. We see that the contributions of the individual intermediate states in the mixing diagram are not intrinsically small – in fact, they begin to contribute at $\mathcal{O}(z^0)$. However, flavour cancellations remove all contributions through $\mathcal{O}(z^2)$ for $\Delta\Gamma_{\rm D}$, so the net result is $\mathcal{O}(z^3)$. Charm mixing clearly experiences a remarkable GIM suppression!

We understand the reason for this. D^0 mixing vanishes in the limit of exact $SU(3)$ flavour symmetry. It is nonzero only because flavour $SU(3)$ is broken, and indeed, D^0 mixing occurs at second order in $SU(3)$ breaking [10]. A factor of z will accompany each order of $SU(3)$ breaking and the rate difference $y_{\rm D}$ will experience an additional factor of z due to helicity suppression.

Of course, this is just the leading order (LO) result in QCD, and we should consider the next-to-leading order result as well,

$$\begin{aligned} x_{\rm D} &= x_{\rm D}^{\rm (LO)} + x_{\rm D}^{\rm (NLO)}, \\ y_{\rm D} &= y_{\rm D}^{\rm (LO)} + y_{\rm D}^{\rm (NLO)}. \end{aligned} \tag{10}$$

Table 2. Results at dimension-six in the OPE.

	LO	NLO	LO + NLO
$y_{\rm D}$	$-(5.7 \to 9.5) \cdot 10^{-8}$	$(3.9 \to 9.1) \cdot 10^{-7}$	$\simeq 6 \cdot 10^{-7}$
$x_{\rm D}$	$-(1.4 \to 2.4) \cdot 10^{-6}$	$(1.7 \to 3.0) \cdot 10^{-6}$	$\simeq 6 \cdot 10^{-7}$

This has been done in [8] and the results are summarized in table 2, which reveals that $y_{\rm D}$ is given by $y_{\rm NLO}$ to a reasonable approximation (due to the removal of helicity suppression by virtual gluons) whereas $x_{\rm D}$ is greatly affected by destructive interference between $x_{\rm LO}$ and $x_{\rm NLO}$. The net effect is to render $y_{\rm D}$ and $x_{\rm D}$ of similar small magnitudes, at least through this order of analysis, as compared to the experiment signal.

It is possible that the quark-level prediction of $x_{\rm D}$ and $y_{\rm D}$ just described might be considerably affected by a higher order in the OPE [11] which suffers less z suppression. Simple dimensional analysis [12] suggests that the magnitudes $x_{\rm D} \sim y_{\rm D} \sim 10^{-3}$ might be achievable, although order-of-magnitude cancellations or enhancements are possible.

2.2 *Hadron-level analysis*

Most of the work involving the hadron degree of freedom has been done on $y_{\rm D}$. One starts with the following general expression for $\Delta\Gamma_{\rm D}$:

$$\Delta\Gamma_{\rm D} = \frac{1}{M_{\rm D}} \,{\rm Im}\, I$$
$$I \equiv \langle \bar{D}^0 | \, i \int {\rm d}^4 x \, T\Big\{ \mathcal{H}_{\rm w}^{|\Delta C|=1}(x) \, \mathcal{H}_{\rm w}^{|\Delta C|=1}(0) \Big\} | D^0 \rangle. \qquad (11)$$

To utilize this relation, one inserts intermediate states between the $|\Delta C| = 1$ weak Hamiltonian densities $\mathcal{H}_{\rm w}^{|\Delta C|=1}$. Although this can be done using either quark or hadron degrees of freedom, let us consider the latter here. Clearly, some knowledge of the matrix elements $\langle n | \mathcal{H}_{\rm w}^{|\Delta C|=1} | D^0 \rangle$ is required.

One approach is to model $|\Delta C| = 1$ decays theoretically and fit the various model parameters to charm decay data. Some time ago, $\Delta\Gamma_{\rm D}$ was determined in this manner and the result $y_{\rm D} \simeq 10^{-3}$ was found [13]. This value is smaller than the recent BaBar and Belle central values.

Alternatively, one can just use the charm decay data and explicitly try to evaluate $\delta\gamma$. The earliest work in this regard focussed on the $P^+P^- = \pi^+\pi^-, K^+K^-, K^-\pi^+, K^+\pi^-$ states [14,15]. In the flavour $SU(3)$ limit, this subset of states gives zero contribution due to cancellations and the data were consistent with a small contribution. A modern version of this approach now exists, although the analysis takes a different slant [10]. They take the charm decay data when available, and for the rest, assume that the only $SU(3)$ breaking is in the phase space. Then, clearly in the decay into 4P states, there will be large $SU(3)$ breaking due to the fact that some states such as 4K states are kinematically not accessible. In [10] it is estimated that these 'kinematically-challenged' sectors can provide enough $SU(3)$ violation to induce $y_{\rm D} \sim 10^{-2}$. This approach is an important advance in our understanding of the subject. At the same time, it is unfortunately more persuasive than compelling due to the uncontrollable uncertainties inherent in this line of reasoning.

The case of the mass difference is somewhat different. To calculate δm, the authors of ref. [16] write a dispersion relation for it in terms of $\delta\Gamma(E)$. But now one needs an extrapolation of $\delta\Gamma(E)$ off-shell and hence a model for this off-shell behaviour. They propose a seemingly reasonable model and estimate δm. Their result is that $x_{\rm D}$ can be of order of (0.1 to 1) times $y_{\rm D}$ but with a sign opposite to $y_{\rm D}$. The final upshot is as follows: for $y_{\rm D}$ the estimate is in the range of 10^{-3} to 10^{-2}; while $x_{\rm D}$ is expected to be in the range $-(10^{-4}$–$10^{-2})$. The relative sign is opposite to the experimental result. The long distance estimate of the mass difference is clearly on a less firm footing than $\delta\Gamma$, the lifetime difference. To summarize the current status of the SM expectations for D mixings; I feel that it is fair to say the following: the width difference ($y_{\rm D}$) is probably correctly accounted by the SM long distance contributions; however, it is quite likely that there are significant new physics contribution to mass difference ($x_{\rm D}$).

3. NP and the width difference

At first glance, it would appear unlikely that NP could affect $y_{\rm D}$ because the particles contributing to the width must be on-shell. Since NP particles will be heavier than the charm mass, 'there can be no NP contribution to $y_{\rm D}$'. This is the conventional wisdom.

However, as explained in ref. [4], NP effects in $\mathcal{H}_{\rm w}^{|\Delta C|=1}$ can generally contribute to $y_{\rm D}$. In the loop amplitude of figure 1, the NP contribution (empirically small for $\Delta C = -1$ processes) arises from either of the two vertices. We represent the NP $\Delta C = -1$ Hamiltonian as (indices i, j, k, ℓ represent colour),

$$\begin{aligned}
&\mathcal{H}_{\rm NP}^{\Delta C=-1} = \sum_{q,q'} D_{qq'}\left[\bar{\mathcal{C}}_1(\mu)\mathcal{O}_1 + \bar{\mathcal{C}}_2(\mu)\mathcal{O}_2\right],\\
&\mathcal{O}_1 = \bar{u}_i\bar{\Gamma}_1 q'_j\ \bar{q}_j\bar{\Gamma}_2 c_i\ ,\\
&\mathcal{O}_2 = \bar{u}_i\bar{\Gamma}_1 q'_i\ \bar{q}_j\bar{\Gamma}_2 c_j,
\end{aligned} \tag{12}$$

where $D_{qq'}$ and the spin matrices $\bar{\Gamma}_{1,2}$ encode the NP model. $\bar{\mathcal{C}}_{1,2}(\mu)$ are Wilson coefficients evaluated at energy scale μ and the flavour sums on q, q' extend over the d, s quarks. This leads to a prediction for the NP contribution to $y_{\rm D}$. For a generic NP interaction, one finds (with the number of colors $N_{\rm c} = 3$)

$$\begin{aligned}
y_{\rm D} = &-\frac{4\sqrt{2}G_{\rm F}}{M_{\rm D}\Gamma_{\rm D}}\sum_{q,q'}\mathbf{V}^*_{cq'}\mathbf{V}_{uq}D_{qq'}\,(K_1\delta_{ik}\delta_{j\ell}\\
&+ K_2\delta_{i\ell}\delta_{jk})\sum_{\alpha=1}^{5} I_\alpha(x,x')\langle\bar{D}^0|\mathcal{O}_\alpha^{ijk\ell}|D^0\rangle,
\end{aligned} \tag{13}$$

Table 3. Some NP models and $y_{\rm D}$.

Model	$y_{\rm D}$	Comment
RPV-SUSY	$6 \cdot 10^{-6}$	Squark exchange
	$-4 \cdot 10^{-6}$	Slepton exchange
Left–right	$-5 \cdot 10^{-6}$	'Manifest'
	$-8.8 \cdot 10^{-5}$	'Nonmanifest'
Multi-Higgs	$2 \cdot 10^{-10}$	Charged Higgs
Extra quarks	10^{-8}	Not little Higgs

where $\{K_\alpha\}$ are combinations of Wilson coefficients,

$$K_1 = \left(\mathcal{C}_1\bar{\mathcal{C}}_1 N_{\rm c} + \left(\mathcal{C}_1\bar{\mathcal{C}}_2 + \bar{\mathcal{C}}_1\mathcal{C}_2\right)\right),$$
$$K_2 = \mathcal{C}_2\bar{\mathcal{C}}_2, \tag{14}$$

and $\{\mathcal{O}_\alpha^{ijk\ell}\}$ are four-quark operators written down in ref. [4]. Numerical results for some NP models are displayed in table 3.

One sees that the entries, aside from R-parity violating SUSY, produce small contributions. We emphasize, however, that eqs (12) and (13) represent general formulae for the contribution of all NP models of $|\Delta C| = 1$ interactions, encompassing those not included in table 3.

4. NP and the mass difference

As the operation of the LHC looms near, the number of potentially viable NP models has never been greater. In this section, the results of [5] are summarized, whose hallmark is the study of many (21 in all) NP models. Perhaps the best way to start is to consider the different ways that 'extras' can be added to the SM:

- Extra gauge bosons (LR models, etc.)
- Extra scalars (multi-Higgs models, etc.)
- Extra fermions (little Higgs models, etc.)
- Extra dimensions (universal extra dims., etc.)
- Extra global symmetries (SUSY, etc.)

Although this approach does not provide a totally clean partition of NP models (e.g. obviously SUSY contains extra particles appearing in other categories), it proved useful to the authors of ref. [5].

The broad menu of NP models which were analysed is listed in table 4. The extensive content of this list (e.g. there are four different SUSY realizations and three involving large extra dimensions) indicates how rich the field of NP models has become.

Suppose a vector-like quark of charge $Q = +2/3$ [17] is added to the SM. Recall that a vector-like quark is one whose electric charge is either $Q = +2/3$ or

Table 4. NP models studied in [5].

Model
Fourth generation
$Q = -1/3$ singlet quark
$Q = +2/3$ singlet quark
Little Higgs
Generic Z'
Family symmetries
Left–right symmetric
Alternate left–right symmetric
Vector leptoquark bosons
Flavour conserving two-Higgs-doublet
Flavour changing neutral Higgs
FC neutral Higgs (Cheng-Sher ansatz)
Scalar leptoquark bosons
Higgsless
Universal extra dimensions
Split fermion
Warped geometries
Minimal supersymmetric standard
Supersymmetric alignment
Supersymmetry with RPV
Split supersymmetry

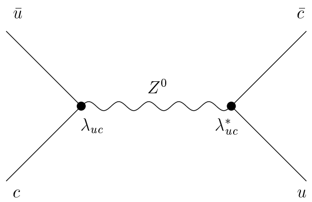

Figure 2. Tree-level contribution from Z^0-exchange.

$Q = -1/3$ and which is an $SU(2)_{\rm L}$ singlet. Both choices of charge are actually well motivated, as such fermions appear explicitly in several NP models. For example, weak isosinglets with $Q = -1/3$ appear in E_6 GUTs [18,19], with one for each of the three generations (D, S and B). Weak iso-singlets with $Q = +2/3$ occur in little Higgs theories [20,21] in which the Standard Model Higgs boson is a pseudo-Goldstone boson, and the heavy isosinglet T quark cancels the quadratic divergences generated by the top quark in the mass of the Higgs boson.

We restrict our attention here to the $Q = +2/3$ case. Since the electroweak quantum number assignments are different from those for the SM fermions, flavour changing neutral current interactions will be generated in the left-handed up-quark sector. Thus, there will also be FCNC couplings with the Z^0 boson [17]. These couplings contain a mixing parameter λ_{uc} which is constrained by the unitarity condition

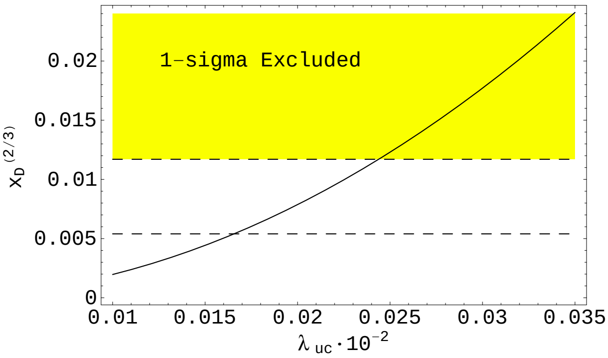

Figure 3. Value of $x_{\rm D}$ as a function of the mixing parameter λ_{uc} in units of 10^{-2} in the $Q = +2/3$ quark singlet model. The 1σ experimental bounds are as indicated, with the shaded area depicting the region that is excluded.

$$\lambda_{uc} \equiv -\left(V^*_{ud}V_{cd} + V^*_{us}V_{cs} + V^*_{ub}V_{cb}\right). \tag{15}$$

A tree-level contribution to $\Delta M_{\rm D}$ is thus generated from Z^0-exchange (see figure 2). It is straightforward to calculate that

$$\begin{aligned} x_{\rm D}^{(2/3)} &= \frac{G_{\rm F}\lambda_{uc}^2}{\sqrt{2}M_{\rm D}\Gamma_{\rm D}} r_1(m_c, M_Z)\langle \bar{D}^0|\mathcal{Q}_1|D^0\rangle \\ &= \frac{2G_{\rm F}f_{\rm D}^2 M_{\rm D}}{3\sqrt{2}\Gamma_{\rm D}} B_{\rm D}\left(\lambda_{uc}\right)^2 r_1(m_c, M_Z), \end{aligned} \tag{16}$$

where we have made use of eqs (4) and (7). The result is displayed in the graph of figure 3, which contrasts a $\pm 1\sigma$ window (dashed lines) about the HFAG central value with the NP prediction $x_{\rm D}^{(2/3)}$ (solid line), as a function of the mixing parameter λ_{uc}. The bound on λ_{uc} from D^0 mixing turns out to be roughly two orders of magnitude better than that from the CKM unitarity constraint.

Upon performing analogous analyses for the other NP models, we arrive at a set of constraints (mainly in the form of graphs) like the one depicted in figure 3. It is not possible here to summarize the results for all 21 models of table 4. One must refer to ref. [5] for that. There are just four models, which are not constrained by the D^0–$\bar{D}^0$ mixing.

It is of interest to briefly consider a few of these, in order to understand how the D^0 mixing constraints can be evaded:

1. Split SUSY [22]: This is a relatively new variant of SUSY (2003–4) in which SUSY breaks at a very large scale $M_S \gg 1000$ TeV. All scalars except the Higgs have mass $M \sim M_S$ whereas fermions have the usual weak-scale mass. It is known that large D^0 mixing in SUSY will involve squark amplitudes. But since squark masses in split SUSY are huge, the mixing becomes suppressed.
2. Universal extra dimensions [23]: UED is a variant of the idea that TeV^{-1}-sized extra dimensions exist. There are no branes appearing in this approach, so all SM fields reside in the bulk and just one extra dimension is usually

considered. Each SM field will have an infinity of KK excitations. It turns out that GIM cancellations suppress all but a few b-quark KK terms, but these are CKM inhibited.

So we see that supressions can arise from more than one source and that the suppressing mechanism will depend on the specific model.

5. Conclusions

At long last, signals for x_{D} and y_{D} have been observed. These experimental findings, although greatly welcome, whet our appetite for ever more precise determinations. Hopefully these will be forthcoming, so we can put aside any lingering concerns that all the excitement has been the result of statistical fluctuations.

The SM analysis is especially difficult in the charm sector. At the quark level, theoretical analysis in the dimension-six sector through NLO gives $x_{\mathrm{D}} \sim y_{\mathrm{D}} \simeq 10^{-6}$. These values are tiny compared to the reported experimental signals. It is evident that the triple sum over the operator dimension d_n, the QCD coupling α_s and the mass expansion parameter z of eq. (9) is slowly convergent. This approach remains inconclusive at best.

A more promising avenue is to study y_{D} with the hadronic degree of freedom. This yields a plausible, and quite possibly correct, explanation for reaching the $y_{\mathrm{D}} \sim 0.01$ level. Again, however, the effect of strong interaction uncertainties mars predictive power. The situation for the mass difference is even murkier, with the possibility that the bulk of the contribution may come from NP.

Finally, the work of refs. [4,5] has explored which NP models can yield sizable values for x_{D}, y_{D} and which cannot. Charm mixing data have been found to infer useful constraints on NP parameter spaces. For future we look forward to smaller errors; and further signals for new physics in rare decay modes of D and in CP violation.

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