

CONSTRAINING NEW PHYSICS FROM $D^0 - \overline{D}^0$ MIXING

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I review constraints on possible New Physics interactions from $D^0 - \overline{D}^0$ mixing measurements. I consider the most general low energy effective Hamiltonian and include leading order QCD running of effective operators. I discuss constraints from an extensive list of popular New Physics models, each of which could be discovered at the LHC, that can generate these operators. In most of the scenarios, strong constraints that surpass those from other search techniques could be placed on the allowed parameter space using the existent evidence for observation of D meson mixing.

1 Introduction

Meson-antimeson mixing has traditionally been of importance because it is sensitive to heavy degrees of freedom that propagate in the underlying mixing amplitudes. Estimates of the charm quark and top quark mass scales were inferred from the observation of mixing in the K^0 and B_d systems, respectively, before these particles were discovered directly.

This success has motivated attempts to indirectly detect New Physics (NP) signals by comparing the observed meson mixing with predictions of the Standard Model (SM). $K^0 - \overline{K}^0$ mixing has historically placed stringent constraints on the parameter space of theories beyond the SM and provides an essential hurdle that must be passed in the construction of models with NP. The large mixing signal in the B_d and B_s systems, observed at the B-factories and the Tevatron collider, can be precisely described in terms of the SM alone, which makes the parameter spaces of various NP models increasingly constrained. These facts influenced theoretical and experimental studies of D^0 flavor oscillations, where the SM mixing rate is sufficiently small that the NP component might be able to compete. There has been a flurry of recent experimental activity

regarding the detection of D^0 - \bar{D}^0 mixing, which marks the first time Flavor Changing Neutral Currents (FCNC) have been observed in the charged $+2/3$ quark sector. With the potential window to discern large NP effects in the charm sector and the anticipated improved accuracy for future mixing measurements, the motivation for a comprehensive up-to-date theoretical analysis of New Physics contributions to D meson mixing is compelling.

The phenomenon of meson-anti-meson mixing occurs in the presence of operators that change quark flavor by two units¹. Those operators can be generated in both the Standard Model and many possible extensions of it. They produce off-diagonal terms in the meson-anti-meson mass matrix, so that the basis of flavor eigenstates no longer coincide with the basis of mass eigenstates. Those two bases, however, are related by a linear transformation,

$$|D_{\frac{1}{2}}\rangle = p|D^0\rangle \pm q|\bar{D}^0\rangle, \quad (1)$$

where the complex parameters p and q are obtained from diagonalizing the D^0 - \bar{D}^0 mass matrix. Neglecting CP-violation leads to $p = q = 1/\sqrt{2}$. The mass and width splittings between those mass eigenstates are given by

$$x_D = \frac{m_1 - m_2}{\Gamma_D}, \quad y_D = \frac{\Gamma_1 - \Gamma_2}{2\Gamma_D}. \quad (2)$$

It is expected that x_D and y_D should be rather small in the Standard Model, which is usually attributed to the absence of superheavy quarks destroying Glashow-Iliopoulos-Maiani (GIM) cancellation. In Eq. (2), Γ_D is the average width of the two neutral D meson mass eigenstates. The quantities which are actually measured in most experimental determinations of the mass and width differences, $y_D^{(\text{CP})}$, x'_D , and y'_D , are defined as

$$\begin{aligned} y_D^{(\text{CP})} &= y_D \cos \phi - x_D \sin \phi \left(\frac{A_m}{2} - A_{\text{prod}} \right) , \\ x'_D &= x_D \cos \delta_{K\pi} + y_D \sin \delta_{K\pi} , \\ y'_D &= y_D \cos \delta_{K\pi} - x_D \sin \delta_{K\pi} , \end{aligned} \quad (3)$$

where $A_{\text{prod}} = (N_{D^0} - N_{\bar{D}^0}) / (N_{D^0} + N_{\bar{D}^0})$ is the so-called production asymmetry of D^0 and \bar{D}^0 (giving the relative weight of D^0 and \bar{D}^0 in the sample) and $\delta_{K\pi}$ is the strong phase difference between the Cabibbo favored and double Cabibbo suppressed amplitudes², which is usually measured in $D \rightarrow K\pi$ transitions. In what follows we shall neglect CP-violating parameters ϕ and A_m . In this limit $y_D^{(\text{CP})} = y_D$. Please see recent reviews^{1,3,4} for more complete analysis.

2 Experimental Constraints on Charm Mixing

The recent interest in D^0 - \bar{D}^0 mixing started with the almost simultaneous observations by the BaBar⁶ and Belle⁷ collaborations of nonzero mixing signals at about the per cent level,

$$y'_D = (0.97 \pm 0.44 \pm 0.31) \cdot 10^{-2} \quad (\text{BaBar}) , \quad (4)$$

$$y_D^{(\text{CP})} = (1.31 \pm 0.32 \pm 0.25) \cdot 10^{-2} \quad (\text{Belle}) . \quad (5)$$

This was soon followed by the announcement by the Belle collaboration of mixing measurements from the Dalitz plot analyses of $D^0 \rightarrow K_S \pi^+ \pi^-$ ⁸,

$$x_D = (0.80 \pm 0.29 \pm 0.17) \cdot 10^{-2} , \quad y_D = (0.33 \pm 0.24 \pm 0.15) \cdot 10^{-2} . \quad (6)$$

A fit to the current database by the Heavy Flavor Averaging Group (HFAG) gives⁵

$$x_D = 9.8^{+2.6}_{-2.7} \cdot 10^{-3} , \quad y_D = (7.5 \pm 1.8) \cdot 10^{-3} , \quad (7)$$

which is obtained assuming no CP-violation affecting mixing. It is important to note that the combined analysis of x_D and y_D excludes the "no-mixing" point $x_D = y_D = 0$ by $6.7\sigma^5$. This fact adds confidence that charm mixing has indeed been observed. Then, a correct interpretation of the results is important. In addition, as with any rare low-energy transition, the question arises on how to use it to probe for physics beyond the Standard Model.

3 Standard Model "background" in $D^0 - \overline{D^0}$ mixing

Theoretical predictions for x_D and y_D obtained in the framework of the Standard Model historically span several orders of magnitude. I will not discuss predictions of the SM for the charm mixing rates here, instead referring the interested reader to recent reviews^{1,3,4}. It might be advantageous to note that there are two approaches to describe $D^0 - \overline{D^0}$ mixing, neither of which give very reliable results because m_c is in some sense intermediate between heavy and light.

The inclusive approach^{10,11} is based on the operator product expansion (OPE). In the formal limit $m_c \gg \Lambda$ limit, where Λ is a scale characteristic of the strong interactions, x_D and y_D can be expanded in terms of matrix elements of local operators. The use of the OPE relies on local quark- hadron duality, and on Λ/m_c being small enough to allow a truncation of the series after the first few terms. This, however, is not realized in charm mixing, as the leading term in $1/m_c$ is suppressed by four and six powers of the strange quark mass for x_D and y_D respectively. The parametrically-suppressed higher order terms in $1/m_c$ can have less powers of m_s , thus being more important numerically¹¹. This results in reshuffling of the OPE series, making it a triple expansion in $1/m_c$, m_s , and α_s . The (numerically) leading term contains over twenty matrix elements of dimension-12, eight-quark operators, which are difficult to compute reliably. A naive power counting then yields $x_D, y_D < 10^{-3}$. The exclusive approach¹² sums over intermediate hadronic states. Since there are cancellations between states within a given $SU(3)$ multiplet, one needs to know the contribution of each state with high precision. However, the D is not light enough that its decays are dominated by a few final states. In the absence of sufficiently precise data, one is forced to use some assumptions. Large effects in y_D appear for decays close to D threshold, where an analytic expansion in $SU(3)_F$ violation is no longer possible. Thus, even though theoretical calculations of x_D and y_D are quite uncertain, the values $x_D \sim y_D \sim 1\%$ are quite natural in the Standard Model¹³.

It then appears that experimental results of Eq. (7) are consistent with the SM predictions. Yet, those predictions are quite uncertain to be subtracted from the experimental data to precisely constrain possible NP contributions. In this situation the following approach can be taken. One can neglect the SM contribution altogether and assume that NP saturates the result reported by experimental collaborations. This way, however, only an upper bound on the NP parameters can be placed. A subtlety of this method of constraining the NP component of the mixing amplitude is related to the fact that the SM and NP contributions can have either the same or opposite signs. While the sign of the SM contribution cannot be calculated reliably due to hadronic uncertainties, x_D computed entirely within a given NP model can be determined rather precisely. This stems from the fact that NP contributions are generated by heavy degrees of freedom making short-distance OPE reliable. This means that only the part of parameter space of NP models that generate x_D of the same sign as observed experimentally can be reliably and unambiguously constrained.

4 New Physics contributions to $D^0 - \overline{D^0}$ mixing

Any NP degree of freedom will generally be associated with a generic heavy mass scale M , at which the NP interaction will be most naturally described. At the scale m_c of the charm

mass, this description will have been modified by the effects of QCD, which should be taken into account. In order to see how NP might affect the mixing amplitude, it is instructive to consider off-diagonal terms in the neutral D mass matrix,

$$\left(M - \frac{i}{2}\Gamma\right)_{12} = \frac{1}{2M_D} \langle \bar{D}^0 | \mathcal{H}_w^{\Delta C=-2} | D^0 \rangle + \frac{1}{2M_D} \sum_n \frac{\langle \bar{D}^0 | \mathcal{H}_w^{\Delta C=-1} | n \rangle \langle n | \mathcal{H}_w^{\Delta C=-1} | D^0 \rangle}{M_D - E_n + i\epsilon} \quad (8)$$

where the first term contains $\mathcal{H}_w^{\Delta C=-2}$, which is an effective $|\Delta C| = 2$ hamiltonian, represented by a set of operators that are local at the $\mu \simeq m_D$ scale. Note that a b -quark also gives a (negligible) contribution to this term. This term only affects x_D , but not y_D .

The second term in Eq. (8) is given by a double insertion of the effective $|\Delta C| = 1$ Hamiltonian $\mathcal{H}_w^{\Delta C=-1}$. This term is believed to give dominant contribution to $D^0 - \bar{D}^0$ mixing in the Standard Model, affecting both x and y . It is generally believed that NP cannot give any sizable contribution to this term, since $\mathcal{H}_w^{\Delta C=-1}$ Hamiltonian also mediates non-leptonic D -decays, which should then also be affected by this NP contribution. I will show that there is a well-defined theoretical limit where NP contribution *dominates* lifetime difference y_D and consider implications of this limit in "real world".

4.1 New Physics in $|\Delta C| = 1$ interactions.

Consider a non-leptonic D^0 decay amplitude, $A[D^0 \rightarrow n]$, which includes a small NP contribution, $A[D^0 \rightarrow n] = A_n^{(\text{SM})} + A_n^{(\text{NP})}$. Here, $A_n^{(\text{NP})}$ is assumed to be smaller than the current experimental uncertainties on those decay rates. This ensures that NP effects cannot be seen in the current experimental analyses of non-leptonic D -decays. One can then write y_D as

$$y_D \simeq \sum_n \frac{\rho_n}{\Gamma_D} A_n^{(\text{SM})} A_n^{(\text{SM})} + 2 \sum_n \frac{\rho_n}{\Gamma_D} A_n^{(\text{NP})} A_n^{(\text{SM})} . \quad (9)$$

The first term of Eq. (schematic) represents the SM contribution to y_D . The SM contribution to y_D is known to vanish in the limit of exact flavor $SU(3)$. Moreover, the first order correction is also absent, so the SM contribution arises only as a *second* order effect¹³. This means that in the flavor $SU(3)$ limit the lifetime difference y_D is dominated by the second term in Eq. (9), i.e. New Physics contributions, even if their contributions are tiny in the individual decay amplitudes¹⁴! A calculation reveals that NP contribution to y_D can be as large as several percent in R-parity-violating SUSY models⁹ or as small as $\sim 10^{-10}$ in the models with interactions mediated by charged Higgs particles¹⁴.

This wide range of theoretical predictions can be explained by two observations. First, many NP affecting $|\Delta C| = 1$ transitions also affect $|\Delta B| = 1$ or $|\Delta S| = 1$ decays or kaon and B-meson mixings, which are tightly constrained. Second, a detailed look at a given NP model that can potentially affect y_D reveals that the NP contribution itself can vanish in the flavor $SU(3)$ limit. For instance, the structure of the NP interaction might simply mimic the one of the SM. Effects like that can occur in some models with extra space dimensions. Also, the chiral structure of a low-energy effective lagrangian in a particular NP model could be such that the leading, mass-independent contribution vanishes exactly, as in a left-right model (LRM). Finally, the NP coupling might explicitly depend on the quark mass, as in a model with multiple Higgs doublets. However, most of these models feature second order $SU(3)$ -breaking already at leading order in the $1/m_c$ expansion. This should be contrasted with the SM, where the leading order is suppressed by six powers of m_s and term of order m_s^2 only appear as a $1/m_c^6$ -order correction.

4.2 New Physics in $|\Delta C| = 2$ interactions.

Though the particles present in models with New Physics may not be produced in charm quark decays, their effects can nonetheless be seen in the form of effective operators generated by the

exchanges of these new particles. Even without specifying the form of these new interactions, we know that their effect is to introduce several $|\Delta C| = 2$ effective operators built out of the SM degrees of freedom.

By integrating out new degrees of freedom associated with new interactions at a scale M , we are left with an effective hamiltonian written in the form of a series of operators of increasing dimension. Operator power counting then tells us the most important contributions are given by the operators of the lowest possible dimension, $d = 6$ in this case. This means that they must contain only quark degrees of freedom and no derivatives. Realizing this, we can write the complete basis of these effective operators, which can be done most conveniently in terms of chiral quark fields,

$$\langle f | \mathcal{H}_{NP} | i \rangle = G \sum_{i=1} C_i(\mu) \langle f | Q_i | i \rangle(\mu) , \quad (10)$$

where the prefactor G has the dimension of inverse-squared mass, the C_i are dimensionless Wilson coefficients, and the Q_i are the effective operators:

$$\begin{aligned} Q_1 &= (\bar{u}_L \gamma_\mu c_L) (\bar{u}_L \gamma^\mu c_L) , & Q_5 &= (\bar{u}_R \sigma_{\mu\nu} c_L) (\bar{u}_R \sigma^{\mu\nu} c_L) , \\ Q_2 &= (\bar{u}_L \gamma_\mu c_L) (\bar{u}_R \gamma^\mu c_R) , & Q_6 &= (\bar{u}_R \gamma_\mu c_R) (\bar{u}_R \gamma^\mu c_R) , \\ Q_3 &= (\bar{u}_L c_R) (\bar{u}_R c_L) , & Q_7 &= (\bar{u}_L c_R) (\bar{u}_L c_R) , \\ Q_4 &= (\bar{u}_R c_L) (\bar{u}_R c_L) , & Q_8 &= (\bar{u}_L \sigma_{\mu\nu} c_R) (\bar{u}_L \sigma^{\mu\nu} c_R) . \end{aligned} \quad (11)$$

In total, there are eight possible operator structures that exhaust the list of possible independent contributions to $|\Delta C| = 2$ transitions. Since these operators are generated at the scale M where the New Physics is integrated out, a non-trivial operator mixing can occur when one takes into account renormalization group running of these operators between the scales M and μ , with μ being the scale where the hadronic matrix elements are computed. We shall work at the renormalization scale $\mu = m_c \simeq 1.3$ GeV. This evolution is determined by solving the RG equations obeyed by the Wilson coefficients,

$$\frac{d}{d \log \mu} \vec{C}(\mu) = \hat{\gamma}^T \vec{C}(\mu) , \quad (12)$$

where $\hat{\gamma}$ represents the matrix of anomalous dimensions of the operators in Eq. (11)¹⁵. Due to the relatively simple structure of $\hat{\gamma}$, one can easily write the evolution of each Wilson coefficient in Eq. (10) from the New Physics scale M down to the hadronic scale μ , taking into account quark thresholds. Corresponding to each of the eight operators $\{Q_i\}$ ($i = 1, \dots, 8$) is an RG factor $r_i(\mu, M)$. The first of these, $r_1(\mu, M)$, is given explicitly by

$$r_1(\mu, M) = \left(\frac{\alpha_s(M)}{\alpha_s(m_t)} \right)^{2/7} \left(\frac{\alpha_s(m_t)}{\alpha_s(m_b)} \right)^{6/23} \left(\frac{\alpha_s(m_b)}{\alpha_s(\mu)} \right)^{6/25} . \quad (13)$$

and the rest can be expressed in terms of $r_1(\mu, M)$ as

$$\begin{aligned} r_2(\mu, M) &= [r_1(\mu, M)]^{1/2} , & r_6(\mu, M) &= r_1(\mu, M) , \\ r_3(\mu, M) &= [r_1(\mu, M)]^{-4} , & r_7(\mu, M) &= r_4(\mu, M) , \\ r_4(\mu, M) &= [r_1(\mu, M)]^{(1+\sqrt{241})/6} , & r_8(\mu, M) &= r_5(\mu, M) , \\ r_5(\mu, M) &= [r_1(\mu, M)]^{(1-\sqrt{241})/6} , \end{aligned} \quad (14)$$

The RG factors are generally only weakly dependent on the NP scale M since it is taken to be larger than the top quark mass, m_t , and the evolution of α_s is slow at these high mass scales. In Table 1, we display numerical values for the $r_i(\mu, M)$ with $M = 1, 2$ TeV and $\mu = m_c \simeq 1.3$ GeV. Here, we compute α_s using the one-loop evolution and matching expressions for perturbative consistency with the RG evolution of the effective hamiltonian. A contribution to $D^0 - \bar{D}^0$

$M(\text{TeV})$	$r_1(m_c, M)$	$r_2(m_c, M)$	$r_3(m_c, M)$	$r_4(m_c, M)$	$r_5(m_c, M)$
1	0.72	0.85	3.7	0.41	2.2
2	0.71	0.84	4.0	0.39	2.3

Table 1: Dependence of the RG factors on the heavy mass scale M .

mixing from a particular NP model can be obtained by calculating matching conditions for the Wilson coefficients C_i at the scale M , running their values down to μ and computing the relevant matrix elements of four-quark operators. A generic model of New Physics would then give the following contribution x_D ,

$$x_D^{NP} = G \frac{f_D^2 B_D m_D}{\Gamma_D} \left[\frac{2}{3} [C_1(m_c) + C_6(m_c)] - \frac{5}{12} [C_4(m_c) + C_7(m_c)] + \frac{7}{12} C_3(m_c) - \frac{5C_2(m_c)}{6} + [C_5(m_c) + C_8(m_c)] \right]. \quad (15)$$

Here we simplified the result by assuming that all non-perturbative ('bag') parameters are equal to $B_D \simeq 0.82$. The Wilson coefficients at the scale μ are related to the Wilson coefficients at the scale M by renormalization group evolution,

$$\begin{aligned} C_1(m_c) &= r_1(m_c, M) C_1(M), \\ C_2(m_c) &= r_2(m_c, M) C_2(M), \\ C_3(m_c) &= \frac{2}{3} [r_2(m_c, M) - r_3(m_c, M)] C_2(M) + r_3(m_c, M) C_3(M), \\ C_4(m_c) &= \frac{8}{\sqrt{241}} [r_5(m_c, M) - r_4(m_c, M)] \left[C_4(M) + \frac{15}{4} C_5(M) \right] \\ &\quad + \frac{1}{2} [r_4(m_c, M) + r_5(m_c, M)] C_4(M), \\ C_5(m_c) &= \frac{1}{8\sqrt{241}} [r_4(m_c, M) - r_5(m_c, M)] [C_4(M) + 64C_5(M)] \\ &\quad + \frac{1}{2} [r_4(m_c, M) + r_5(m_c, M)] C_5(M), \\ C_6(m_c) &= r_6(m_c, M) C_6(M), \\ C_7(m_c) &= \frac{8}{\sqrt{241}} [r_8(m_c, M) - r_7(m_c, M)] \left[C_7(M) + \frac{15}{4} C_8(M) \right] \\ &\quad + \frac{1}{2} [r_7(m_c, M) + r_8(m_c, M)] C_7(M), \\ C_8(m_c) &= \frac{1}{8\sqrt{241}} [r_7(m_c, M) - r_8(m_c, M)] [C_7(M) + 64C_8(M)] \\ &\quad + \frac{1}{2} [r_7(m_c, M) + r_8(m_c, M)] C_8(M), \end{aligned} \quad (16)$$

A contribution of each particular NP model can then be studied using Eq. (15). Even before performing such an analysis, one can get some idea what energy scales can be probed by $D^0 - \bar{D}^0$ mixing. Setting $G = 1/M^2$ and $C_i(M) = 1$, we obtain $M \sim 10^3$ TeV. More realistic models can be probed in the region of several TeV, which is very relevant for LHC phenomenology applications.

A program described above has been recently executed¹⁵ for 21 well-motivated NP models, which will be actively studied at LHC. The results are presented in Table 2. As can be seen, out of 21 models considered, only four received no useful constraints from $D^0 - \bar{D}^0$ mixing. More informative exclusion plots can be found in that paper¹⁵ as well. It is interesting to note that

Model	Approximate Constraint
Fourth Generation	$ V_{ub'}V_{cb'} \cdot m_{b'} < 0.5$ (GeV)
$Q = -1/3$ Singlet Quark	$s_2 \cdot m_S < 0.27$ (GeV)
$Q = +2/3$ Singlet Quark	$ \lambda_{uc} < 2.4 \cdot 10^{-4}$
Little Higgs	Tree: See entry for $Q = -1/3$ Singlet Quark Box: Parameter space can reach observed x_D
Generic Z'	$M_{Z'}/C > 2.2 \cdot 10^3$ TeV
Family Symmetries	$m_1/f > 1.2 \cdot 10^3$ TeV (with $m_1/m_2 = 0.5$)
Left-Right Symmetric	No constraint
Alternate Left-Right Symmetric	$M_R > 1.2$ TeV ($m_{D_1} = 0.5$ TeV) $(\Delta m/m_{D_1})/M_R > 0.4$ TeV $^{-1}$
Vector Leptoquark Bosons	$M_{VLQ} > 55(\lambda_{PP}/0.1)$ TeV
Flavor Conserving Two-Higgs-Doublet	No constraint
Flavor Changing Neutral Higgs	$m_H/C > 2.4 \cdot 10^3$ TeV
FC Neutral Higgs (Cheng-Sher)	$m_H/ \Delta_{uc} > 600$ GeV
Scalar Leptoquark Bosons	See entry for RPV SUSY
Higgsless	$M > 100$ TeV
Universal Extra Dimensions	No constraint
Split Fermion	$M/ \Delta y > (6 \cdot 10^2$ GeV)
Warped Geometries	$M_1 > 3.5$ TeV
MSSM	$ (\delta_{12}^u)_{\text{LR},\text{RL}} < 3.5 \cdot 10^{-2}$ for $\tilde{m} \sim 1$ TeV $ (\delta_{12}^u)_{\text{LL},\text{RR}} < .25$ for $\tilde{m} \sim 1$ TeV
SUSY Alignment	$\tilde{m} > 2$ TeV
Supersymmetry with RPV	$\lambda'_{12k}\lambda'_{11k}/m_{\tilde{d}_{R,k}} < 1.8 \cdot 10^{-3}/100$ GeV
Split Supersymmetry	No constraint

Table 2: Approximate constraints on NP models from D^0 mixing.

some models *require* large signals in the charm system if mixing and FCNCs in the strange and beauty systems are to be small (as in, for example, the SUSY alignment model^{16,17}).

5 Conclusions

I reviewed implications of recent measurement of $D^0 - \bar{D}^0$ mixing rates for constraining models of New Physics. A majority of considered models received competitive constraints from $D^0 - \bar{D}^0$ mixing measurements despite hadronic uncertainties that plague SM contributions. It should be noted that vast majority of predictions of NP models do not suffer from this uncertainty, and can be computed reliably, if lattice QCD community provides calculations of matrix elements of four-fermion operators Eq. (11).

Another possible manifestation of new physics interactions in the charm system is associated with the observation of (large) CP-violation^{1,4,18}. This is due to the fact that all quarks that build up the hadronic states in weak decays of charm mesons belong to the first two generations. Since 2×2 Cabibbo quark mixing matrix is real, no CP-violation is possible in the dominant tree-level diagrams which describe the decay amplitudes. CP-violating amplitudes can be introduced in the Standard Model by including penguin or box operators induced by virtual b -quarks. However, their contributions are strongly suppressed by the small combination of CKM matrix elements $V_{cb}V_{ub}^*$. It is thus widely believed that the observation of (large) CP violation in charm decays or mixing would be an unambiguous sign for New Physics.

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