New Physics contributions to the lifetime difference in $D^0$-$\bar{D}^0$ mixing

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The first general analysis of New Physics contributions to the $D^0$-$\bar{D}^0$ lifetime difference (equivalently $\Delta\Gamma_D$) is presented. The extent to which New Physics (NP) contributions to $|\Delta C| = 1$ processes can produce effects in $\Delta\Gamma_D$, even if such NP contributions are undetectable in the current round of $D^0$ decay experiments, is studied. New Physics models which do and do not dominate the lifetime difference in the flavor $SU(3)$ limit are identified. Specific examples are provided.

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Quantum mechanical meson-antimeson oscillations are sensitive to heavy degrees of freedom which propagate in the underlying mixing amplitudes. The observation of mixing in the $K^0$ and $B_d$ systems thus implied the existence respectively of the charm and top quarks before these particles were discovered. In like manner, by comparing observed meson mixing with predictions of the Standard Model (SM) modern experimental studies have been able to constrain models of New Physics (NP).

Which system of mixed mesons is likely to produce evidence for NP? It has become clear from B-factories and B_d systems thus implied the existence respectively of the charm and top quarks before these particles were discovered. In like manner, by comparing observed meson mixing with predictions of the Standard Model (SM) modern experimental studies have been able to constrain models of New Physics (NP).

The most natural place for NP to affect mixing amplitudes is in the charmless intermediate state $n$. In this Letter, we show that this is not necessarily so and consider several NP models to illustrate our point.

Consider a $D^0$ decay amplitude which includes a small charmless intermediate state $n$. This relation shows that $\Delta\Gamma_D$ is driven by transitions $D^0, \bar{D}^0 \to n$, i.e. physics of the $|\Delta C| = 1$ sector. It turns out that experimentally observed $D^0$ decays agree reasonably well with SM estimates. To date, no clear signals of NP have been observed. As such, it is currently accepted that $\Delta\Gamma_D$ is dominated by the SM contribution. In this Letter, we show that this is not necessarily so and consider several NP models to illustrate our point.

Let us introduce standard notation for $\Delta\Gamma_D$ and $\Delta M_D$ by employing the dimensionless forms,

$$y = \frac{\Delta\Gamma_D}{2\Gamma_D}, \quad x = \frac{\Delta M_D}{\Gamma_D}. \tag{2}$$

Given CP-conservation, we can express $y$ as an absorptive part of Eq. (1),

$$y = \frac{1}{\Gamma_D} \sum_n \rho_n \langle \bar{D}^0 | H_w^{\Delta C = -1} | n \rangle \langle n | H_w^{\Delta C = -1} | D^0 \rangle, \tag{3}$$

where $\rho_n$ is the phase space function that corresponds to charmless intermediate state $n$. This relation shows that $\Delta\Gamma_D$ is driven by transitions $D^0, \bar{D}^0 \to n$, i.e. physics of the $|\Delta C| = 1$ sector. It turns out that experimentally observed $D^0$ decays agree reasonably well with SM estimates. To date, no clear signals of NP have been observed. As such, it is currently accepted that $\Delta\Gamma_D$ is dominated by the SM contribution. In this Letter, we show that this is not necessarily so and consider several NP models to illustrate our point.

Consider a $D^0$ decay amplitude which includes a small NP contribution, $A[D^0 \to n] = A_n^{(SM)} + A_n^{(NP)}$. Here, $A_n^{(NP)}$ is assumed to be smaller than the current experimental uncertainties on those decay rates. Then it is a good approximation to write Eq. (3) in the form

$$y \approx \sum_n \frac{\rho_n}{\Gamma_D} A_n^{(SM)} + 2 \sum_n \frac{\rho_n}{\Gamma_D} A_n^{(NP)} A_n^{(SM)} \tag{4}$$

The first term in this equation corresponds to SM interactions at both vertices in Fig. 1, whereas for the second term, there is one SM vertex and one NP vertex.

The SM contribution to $y$ is known to vanish in the limit of exact flavor $SU(3)$ symmetry. Moreover as was shown in [3], the first order correction is also absent, so the SM contribution arises only as a second order effect. Thus, those NP contributions which do not vanish in the flavor $SU(3)$ limit must determine the lifetime difference there,

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

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even if their contributions are tiny in the individual decay amplitudes. The same reasoning can be applied to $x$ since a dispersion relation relates $x$ to $y$.

Of course, flavor $SU(3)$ symmetry is broken in the real world. Just how large this effect is on $D^0$-$\bar{D}^0$ mixing in the SM is controversial, with estimates for $y$ ranging from a percent to orders of magnitude smaller. The current experimental bounds on $y$ and $x$ are

$$y < 0.008 \pm 0.005, \quad x < 0.029 \text{ (95\% C.L.)}.$$ 

A NP $|\Delta C| = 1$ interaction can have a measureable effect on the value of $y$ (and of $x$) if the true SM value for $y$ does not near the top of the range of predictions. We shall assume that this is the case.

Using the completeness relation and Eq. (3), the NP contribution to the $D^0$-$\bar{D}^0$ lifetime difference becomes

$$y = \frac{2}{M_d\Gamma_D} \langle D^0 | \text{Im} \mathcal{T} | D^0 \rangle,$$ 

$$\mathcal{T} = i \int d^4x T \left( \mathcal{H}^{AC=1}_{SM} - \mathcal{H}^{AC=1}_{NP} \right).$$

We represent the NP $\Delta C = -1$ hamiltonian as

$$\mathcal{H}^{AC=1}_{NP} = \sum_{q,q'} D_{qq'} \left[ \Gamma_1(q_1) + \Gamma_2(q_2) \right],$$

where the spin matrices $\Gamma_{1,2}$ can have arbitrary Dirac structure, $\Gamma_1,2(q)$ are Wilson coefficients evaluated at energy scale $\mu$ and the flavor sums on $q,q'$ extend over the $d,s$ quarks. We shall expand the time-ordered product in an operator product expansion (OPE), i.e. in terms of local operators of increasing dimension.

The leading term in the OPE is simply that depicted in Fig. 4. For a generic NP interaction, we calculate that

$$y = -\frac{4\sqrt{2}G_F}{M_d\Gamma_D} \sum_{q,q'} \bar{V}_{eq} V_{uq'} D_{qq'} (K_1 \delta_{jk} \delta_{j\ell})$$

$$+ K_2 \delta(k_1) \sum_{\alpha \in 1} I_\alpha(x,x') \langle D^0 | O_\alpha^{ijkl} | D^0 \rangle,$$ 

with the number of colors $N_c = 3$. The operators $O_\alpha^{ijkl}$ in Eq. (7) are defined as

$$O_\alpha^{ijkl} = \bar{q}_k \Gamma_\mu \bar{p}_e \Gamma_2 c_j \bar{p}_e \Gamma_1 \Gamma_\nu \Gamma_\mu c_i.$$

Examples of New Physics Models

In what follows, we distinguish between NP models which vanish in the limit of $SU(3)$ flavor symmetry from those which do not.

Nonzero $SU(3)$ Limit: For NP models with flavor-dependent couplings $D_{qq'}$, it is possible to obtain contributions that are nonzero in the flavor $SU(3)$ limit. For
instance, the above type which extends the SM by including new
$tors$ $\lambda_{qq'} = \lambda_{qq'}/\Lambda^2$, where $\Lambda$ is the NP
mass scale. We find
\[ y_{\text{VLH}} = \frac{C\lambda}{\Lambda^2} \left[ \frac{K_1 + 2K_2}{2} \langle Q \rangle + (K_2 - K_1) \langle Q_S \rangle \right], \] (12)
where $C \equiv \sqrt{2}G_F m_W^2/(3\pi M_D \Gamma_D)$ and
\[ \bar{\lambda} = \lambda_{sd} - \lambda (\lambda_{dd} - \lambda_{ss}) - \lambda^2 (\lambda_{ds} + \lambda_{sd}) \] (13)
is the combination of NP couplings to the $s, d$ quarks.

It follows from Eq. (13) that if all NP couplings $\lambda_{qq'}$
of the same order, $y_{\text{VLH}}$ is nonzero in the flavor $SU(3)$
limit. The result for a penguin-like NP contribution $y_{\text{VLH}}$ can be
obtained if one sets $\lambda_{dd} = \lambda_{ss} = 0$. In this case the
same conclusion holds if $\lambda_{dd} \neq \lambda_{ss}$ (which however
is not easy to arrange) unless a (generally tiny) $\langle A_{\text{NP}}^\mu \rangle$
term is also included in Eq. (11). In what follows, we will be
neglecting QCD running of the local operators generated
by the NP interaction.}

Models with extra vector-like quarks: Consider a model
of the above type which extends the SM by including
singlet quarks in a vector-like representation \([10]\).
In this instance, the $Z$-boson has additional flavor-changing
couplings. For example, assume both up-type and down-
type exotic quarks $U_{a,i}$, $D_{a,i}$ are present (indices $a, i$
both flavor and color respectively). Then the flavor-
changing couplings are described by
\[ \mathcal{L}_{xQ} = -\frac{g}{2\cos\theta_W} J_{\mu}^{\text{(NC)}} Z^\mu \text{ h.c.}, \] (14)
\[ J_{\mu}^{\text{(NC)}} = U_{ab}^{(u)} \gamma_{\mu} U_{b,i} + U_{ab}^{(d)} \gamma_{\mu} D_{b,i}, \]
with flavor-changing couplings in both up and down
sectors. The lifetime difference for this model can be
obtained from Eq. (12) by substituting $\lambda_{sd} = U_{cu}^{(u)} U_{sd}^{(d)}$,
$\lambda_{ds} = \lambda_{dd} = \lambda_{ss} = 0$ and $\Lambda = \sqrt{2}/G_F$. This model
is well-constrained from measurements of the mass
differences in $K \overline{K}$ and $D \overline{D}$ mixing. For $U_{ab}^{(u)} \sim 10^{-3}$ and
$U_{ab}^{(d)} \sim 10^{-4}$ we get $y \sim 10^{-8}$, of the same order of
magnitude as $\g_{\text{SM}}^{\text{(LO)}}$. Above it is worth noting that the little
Higgs models, which have similar low-energy signatures,
are not constrained by the measurements of lifetime
difference, as they do not have flavor-changing couplings
for the down quark sector (flavor-conserving contributions
cancel out in $y$).

**SUSY without R-parity (slepton exchange):** Another
example of a contribution which survives $\g_{\text{SM}}^{\text{(LO)}}$. In this
model, there are flavor-changing interactions of sleptons
that can be obtained from the lagrangian
\[ \mathcal{L}_R = \lambda_{ij\ell} L_{i\ell} Q_j D_\ell^c, \] (15)
as well as the interactions mediated by squarks discussed
below. The slepton-mediated interaction is not sup-
pressed in the flavor $SU(3)$ limit and leads to
\[ y_R = \frac{C'\bar{\lambda} M_F}{m_F^2} \left[ \langle C_2 - 2C_1 \rangle \langle Q' \rangle + \langle C_1 - 2C_2 \rangle \langle Q' \rangle \right], \] (16)
where $C' = -G_F m_F^2/(6\sqrt{2}\pi M_D \Gamma_D)$, $M_F$ is a slepton
mass, $\bar{\lambda}$ is given by Eq. (13) with $\lambda_{sd} = \lambda_{112} \lambda_{212} \leq 1 \times 10^{-5}$, $\lambda_{ss} = \lambda_{111} \lambda_{211} \leq 5 \times 10^{-5}$, $\lambda_{dd} = \lambda_{121} \lambda_{222} \leq 5 \times 10^{-5}$, $\lambda_{ds} = \lambda_{111} \lambda_{222} \leq 5 \times 10^{-2}$ \([11]\), and $\langle Q' \rangle$ is
\[ \langle Q' \rangle = \langle D^\dagger \pi_{i\mu} P_L c_i \pi_{j\nu} \gamma^\mu P_R \ell_j D^0 \rangle. \] (17)
Operators with a tilde are obtained by swapping color
indices in the charm quark operators. Using factorization
to estimate matrix elements of the above operators and
taking for definiteness $M_F = 100$ GeV, we arrive at
$y_R \approx -3.7\%$. The contribution due to squark exchange
vanishes in the flavor $SU(3)$ limit and is given below.

**Zero SU(3) Limit:** There are several reasons that
some NP models vanish in the flavor $SU(3)$ limit. First,
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. Second, 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). Third, the NP coupling might
explicitly depend on the quark mass, as in a model with
multiple Higgs doublets. There, the charged Higgs
couplings explicitly depend on quark mass. 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 the second order only
appears as a $1/m_c^6$-order correction.

**Left-right models:** Left-right models (LRM) provide
new tree-level contributions mediated by right-handed
($W^{(R)}$) bosons \([12]\). The relevant effective lagrangian is
\[ \mathcal{L}_LR = -\frac{g_{LR}}{\sqrt{2}} V_{ik}^{(R)} \pi_{a,i} \gamma^\mu P_R d_{b,i} W_{\mu}^{(R)} + \text{h.c.}, \] (18)
where $\pi_{i\mu}$ are the coefficients of the right-handed CKM
matrix. This leads to a local $\Delta C = -1$ hamiltonian as
in Eq. (6) with $\Gamma_1 = \Gamma_2 = \gamma_{15} P_R$. Since current experi-
mental limits allow $W^{(R)}$ masses as low as a TeV \([3]\), a
sizable contribution to $y$ is quite possible. Using Eq. (18),
we obtain
\[ y_{LR} = -C_{LR} V_{cs}^{(R)} V_{us}^{(R)*} \left[ \Gamma_1 \langle Q' \rangle + \Gamma_2 \langle Q' \rangle \right], \] (19)
where $C_{L,R} \equiv \mathcal{L}^{(R)}_{F} G_{F} m_{q}^{2}/(\pi M_{D} \Gamma_{D})$, $G_{F}^{(R)}/\sqrt{2} \equiv g_{u}^{2}/8M_{W}^{2}$, $C_{1,2}$ are the SM Wilson coefficients and the operators appearing in Eq. (19) are given in Eq. (17). Using (3), we obtain numerical values for two possible realizations: (i) `Manifest LR' ($V^{(L)} = V^{(R)}$) gives $y_{L,R} = -8.4 \cdot 10^{-5}$ with $M_{W} = 1.6$ TeV and (ii) `Non-manifest LR' ($V_{ij}^{(R)} \sim 1$) gives $y_{L,R} = -8.8 \cdot 10^{-5}$ with $M_{W} = 0.8$ TeV. In both cases we take $g_{R} = g_{L}$.

**Multi-Higgs models**: A popular realization of this type is the two Higgs doublet model (2HDM) with natural flavor conservation. This model presents new tree-level contributions mediated by charged Higgs bosons and leads to the local four fermion interaction $\mathcal{H}_{\rm ChH}$:

$$\mathcal{H}_{\rm ChH} = -\frac{\sqrt{2}G_{F}}{M_{H}^{2}} \sum_{i} q_{i}^{T} q_{i}^{T} c_{j}^{T} c_{j}^{T},$$

(20)

where the vertices $q_{i}$ are given earlier. Using values $M_{H} = 85$ GeV and $\cot \beta = 0.55$ consistent with constraints obtained from (3), we obtain $y_{q_{i}}^{\rm ChH} \approx 2 \cdot 10^{-10}$.

**SU(3) suppression of the SM amplitude**

This baryon-number violating squark exchanges arise from the lagrangian $\mathcal{L}_{\rm R}$:

$$\mathcal{L}_{\rm R} = \lambda_{ij}^{\nu} U_{ij}^{\nu} D_{j}^{\nu} D_{j}^{\nu}.$$

(23)

This interaction has the same Dirac structure as the LRM discussed earlier and leads to

$$y_{R}^{\nu} = -x_{s} C_{ij}^{\nu} \frac{\lambda_{ij}^{\nu} \lambda_{2k}^{\nu}}{M_{\nu}^{2}} \left[ C_{2}(Q^{\nu}) + C_{1} N_{\nu}(Q^{\nu}) \right],$$

(24)

where $C_{ij}^{\nu} = G_{F} \lambda/(2m_{D} \Gamma_{D})$, $M_{\nu}$ is a squark mass and the matrix elements $\langle Q^{\nu} \rangle$, $\langle \bar{Q}^{\nu} \rangle$ are given earlier.

Using factorization for the matrix elements, $\lambda_{2k}^{\nu} \lambda_{2k}^{\nu} \sim 3 \cdot 10^{-4}$, and taking $M_{\nu} = 100$ GeV, we arrive at the result $y_{R}^{\nu} \approx 6.4 \cdot 10^{-6}$.

In conclusion, we have explored how NP contributions can influence the lifetime difference $\Delta$ in the charm system. We argued that the NP signal is dominant in the formal flavor $SU(3)$ limit. We also showed that, for some NP models, it is possible that small NP contributions to $\Delta C = 1$ processes produce substantial effects in the $\Delta^{0}$ lifetime difference, even if such contributions are currently undetectable in the experimental analyses of charmed meson decays. Coupled with a known difficulty in computing SM contributions to D-meson decay amplitudes, it might be advantageous to use experimental constraints on $y$ in order to test various NP scenarios due to better theoretical control over the NP contribution and SU(3) suppression of the SM amplitude.

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