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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 the Tevatron collider that hopes for spectacular NP contributions to B_d and B_s oscillations have come to naught. The large SM mixing successfully describes all available experimental data. The only flavor oscillation not yet observed is that of the charmed meson D^0 , where SM mixing is very small and the NP component can stand out. [1]

Charm mixing arises from $|\Delta C| = 2$ interactions that generate off-diagonal terms in the neutral D mass matrix. To second order, the D^0 -to- \bar{D}^0 matrix element is

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

$$+ \frac{1}{2M_D} \sum_n \frac{\langle \bar{D}^0 | \mathcal{H}_w^{\Delta C=-1} | n \rangle \langle n | \mathcal{H}_w^{\Delta C=-1} | D^0 \rangle}{M_D - E_n + i\epsilon}$$

where $\mathcal{H}_w^{|\Delta C|=1,2}$ is the effective $|\Delta C| = 1, 2$ hamiltonian. The most natural place for NP to affect mixing amplitudes is in the $|\Delta C| = 2$ piece, which corresponds to a local interaction at the charm quark mass scale. This local interaction cannot, however, affect $\Delta\Gamma_D$ because it does not have an absorptive part.

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

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

Given CP-conservation, we can express y as an absorptive

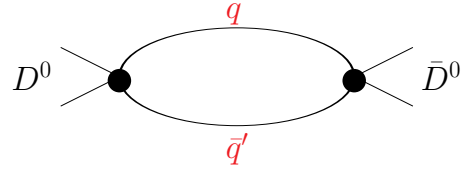


FIG. 1: Loop diagram for $D^0 \rightarrow \bar{D}^0$.

part of Eq. (1),

$$y = \frac{1}{\Gamma_D} \sum_n \rho_n \langle \bar{D}^0 | \mathcal{H}_w^{\Delta C=-1} | n \rangle \langle n | \mathcal{H}_w^{\Delta C=-1} | D^0 \rangle, \quad (3)$$

where ρ_n is a phase space function that corresponds to charmless intermediate state n . This relation shows that $\Delta\Gamma_D$ is driven by transitions $D^0, \bar{D}^0 \rightarrow n$, *i.e.* physics of the $|\Delta C| = 1$ sector. It turns out that experimentally observed D^0 decays agree reasonably well with SM estimates [2]. To date, no clear signals of NP have been observed [4]. 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, $A[D^0 \rightarrow n] = A_n^{(\text{SM})} + A_n^{(\text{NP})}$. Here, $A_n^{(\text{NP})}$ is assumed to be smaller than the current experimental uncertainties on those decay rates. Then it is a good approximation to write Eq. (3) in the form

$$y \simeq \sum_n \frac{\rho_n}{\Gamma_D} A_n^{(\text{SM})} \bar{A}_n^{(\text{SM})} + 2 \sum_n \frac{\rho_n}{\Gamma_D} A_n^{(\text{NP})} \bar{A}_n^{(\text{SM})}. \quad (4)$$

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 $SU(3)$ [5]. Moreover as was shown in [6], 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,

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 [7].

Of course, flavor $SU(3)$ symmetry is broken in the real world. Just how large this effect is on $D^0\text{-}\bar{D}^0$ mixing in the SM is controversial, with estimates for y ranging from a percent [7] to orders of magnitude smaller [2, 8]. The current experimental bounds on y and x are [3]

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

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 $D^0\text{-}\bar{D}^0$ lifetime difference becomes

$$y = \frac{2}{M_D \Gamma_D} \langle \bar{D}^0 | \text{Im } \mathcal{T} | D^0 \rangle, \quad (5)$$

$$\mathcal{T} = i \int d^4x T (\mathcal{H}_{SM}^{\Delta C=-1}(x) \mathcal{H}_{NP}^{\Delta C=-1}(0)) .$$

We represent the NP $\Delta C = -1$ hamiltonian as¹

$$\mathcal{H}_{NP}^{\Delta C=-1} = \sum_{q,q'} D_{qq'} [\bar{\mathcal{C}}_1(\mu) Q_1 + \bar{\mathcal{C}}_2(\mu) Q_2], \quad (6)$$

$$Q_1 = \bar{u}_i \bar{\Gamma}_1 q'_j \bar{q}_j \bar{\Gamma}_2 c_i, \quad Q_2 = \bar{u}_i \bar{\Gamma}_1 q'_i \bar{q}_j \bar{\Gamma}_2 c_j,$$

where the spin matrices $\bar{\Gamma}_{1,2}$ can have arbitrary Dirac structure, $\bar{\mathcal{C}}_{1,2}(\mu)$ are Wilson coefficients evaluated at energy scale μ and the flavor sums on q, q' extend over the d, s quarks. We shall expand the time-ordered product of Eq. (5) in an operator product expansion (OPE), *i.e.* in terms of local operators of increasing dimension [9].

The leading term in the OPE is simply that depicted in Fig. 1. For a generic NP interaction, we calculate that

$$y = - \frac{4\sqrt{2}G_F}{M_D \Gamma_D} \sum_{q,q'} \mathbf{V}_{cq'}^* \mathbf{V}_{uq} D_{qq'} (K_1 \delta_{ik} \delta_{j\ell} + K_2 \delta_{i\ell} \delta_{jk}) \sum_{\alpha=1}^5 I_\alpha(x, x') \langle \bar{D}^0 | \mathcal{O}_\alpha^{ijkl} | D^0 \rangle, \quad (7)$$

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

$$K_1 = (\mathcal{C}_1 \bar{\mathcal{C}}_1 N_c + (\mathcal{C}_1 \bar{\mathcal{C}}_2 + \bar{\mathcal{C}}_1 \mathcal{C}_2)), \quad K_2 = \mathcal{C}_2 \bar{\mathcal{C}}_2 \quad (8)$$

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

$$\mathcal{O}_1^{ijkl} = \bar{u}_k \Gamma_\mu \gamma_\nu \bar{\Gamma}_2 c_j \bar{u}_\ell \bar{\Gamma}_1 \gamma^\nu \Gamma^\mu c_i$$

$$\begin{aligned} \mathcal{O}_2^{ijkl} &= \bar{u}_k \Gamma_\mu \not{p}_c \bar{\Gamma}_2 c_j \bar{u}_\ell \bar{\Gamma}_1 \not{p}_c \Gamma^\mu c_i \\ \mathcal{O}_3^{ijkl} &= \bar{u}_k \Gamma_\mu \bar{\Gamma}_2 c_j \bar{u}_\ell \bar{\Gamma}_1 \not{p}_c \Gamma^\mu c_i \\ \mathcal{O}_4^{ijkl} &= \bar{u}_k \Gamma_\mu \not{p}_c \bar{\Gamma}_2 c_j \bar{u}_\ell \bar{\Gamma}_1 \Gamma^\mu c_i \\ \mathcal{O}_5^{ijkl} &= \bar{u}_k \Gamma_\mu \bar{\Gamma}_2 c_j \bar{u}_\ell \bar{\Gamma}_1 \Gamma^\mu c_i, \end{aligned} \quad (9)$$

where p_c is the charm-quark momentum operator, $\Gamma_\mu \equiv \gamma_\mu P_L$, $P_L \equiv (1 + \gamma_5)/2$ and later we also use $P_R \equiv (1 - \gamma_5)/2$. The coefficients $I_\alpha(x, x')$ in Eq. (7) are

$$\begin{aligned} I_1(x, x') &= -\frac{k^* m_c}{48\pi} [1 - 2(x + x') + (x - x')^2] \\ I_2(x, x') &= -\frac{k^*}{24\pi m_c} [1 + (x + x') - 2(x - x')^2] \\ I_3(x, x') &= \frac{k^*}{8\pi} \sqrt{x} (1 + x' - x) \\ I_4(x, x') &= -\frac{k^*}{8\pi} \sqrt{x'} (1 - x' + x) \\ I_5(x, x') &= \frac{k^* m_c}{4\pi} \sqrt{xx'} \end{aligned} \quad (10)$$

where $k^* \equiv (m_c/2)[1 - 2(x + x') + (x - x')^2]^{1/2}$ with $x \equiv m_q^2/m_c^2$ and $x' \equiv m_{q'}^2/m_c^2$.

Equations (7)-(10) represent the basic formulas in our analysis. Hereafter, we take $m_d = 0$ and express results in terms of the Wolfenstein parameter $\lambda \equiv \mathbf{V}_{us} = -\mathbf{V}_{cd} \simeq 0.22$ and $x_s \equiv m_s^2/m_c^2 \simeq 0.006$. All our results are presented to leading order in x_s . Finally, predictions using NP masses and couplings other than the ones assumed here can be obtained via simple scaling.

The SM hamiltonian $\mathcal{H}_{SM}^{\Delta C=-1}$ is recovered in Eq. (6) by setting $D_{qq'} = -(G_F/\sqrt{2}) \mathbf{V}_{cq}^* \mathbf{V}_{uq'}$, $\bar{\mathcal{C}}_i \rightarrow \mathcal{C}_i$, $\bar{\Gamma}_{1,2} \rightarrow \Gamma_\mu$, and the known SM result $y_{SM}^{(LO)}$ easily follows,

$$\begin{aligned} y_{SM}^{(LO)} &= \frac{G_F^2 m_c^2 \lambda^2 x_s^3}{2\pi M_D \Gamma_D} (K_2 - K_1) [\langle Q \rangle + 4\langle Q_S \rangle], \\ \langle Q \rangle &= \langle \bar{D}^0 | \bar{u}_i \Gamma_\mu c_i \bar{u}_j \Gamma^\mu c_j | D^0 \rangle, \\ \langle Q_S \rangle &= \langle \bar{D}^0 | \bar{u}_i P_R c_i \bar{u}_j P_R c_j | D^0 \rangle. \end{aligned} \quad (11)$$

Note that this contribution is suppressed by six powers of m_s and is therefore tiny, $y_{SM}^{(LO)} \sim 10^{-8}$ [1]. This is because the GIM mechanism requires four chirality flips (strange quark mass insertions) for the intermediate quarks plus additional helicity flip due to the pseudoscalar initial state. [6].

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

¹ Throughout this paper, we reserve indices i, j, k, ℓ for color.

these, the main contribution will come from the operators $\mathcal{O}_{1,2}$ (as $\mathcal{O}_{3,4,5}$ are suppressed by powers of m_s/m_c). The two most common scenarios involve (V-A) \otimes (V-A) and (S-P) \otimes (S+P) couplings.

As an example, consider a NP model whose low energy effective hamiltonian is represented by a four-fermion operator with vectorial left-handed interactions, *i.e.* $\bar{\Gamma}_1 = \bar{\Gamma}_2 = \gamma_\mu P_L$ and $D_{qq'} = \lambda_{qq'}/\Lambda^2$, where Λ is the NP mass scale. We find

$$y_{\text{VLH}} = \frac{C\tilde{\lambda}}{\Lambda^2} \left[\frac{K_1 + 2K_2}{2} \langle Q \rangle + (K_2 - K_1) \langle Q_S \rangle \right], \quad (12)$$

where $C \equiv \sqrt{2}G_F m_c^2 / (3\pi M_D \Gamma_D)$ and

$$\tilde{\lambda} = \lambda_{sd} - \lambda(\lambda_{dd} - \lambda_{ss}) - \lambda^2(\lambda_{ds} + \lambda_{sd}) \quad (13)$$

is the combination of NP couplings to the s, d quarks.

It follows from Eq. (13) that if all NP couplings $\lambda_{qq'}$ are of the same order, y_{VLH} is nonzero in the flavor $SU(3)$ limit. The result for a penguin-like NP contribution can be obtained if one sets $\lambda_{sd} = \lambda_{ds} = 0$. In this case the same conclusion holds if $\lambda_{dd} \neq \lambda_{ss}$ (which is however not easy to arrange) unless a (generally tiny) $|A_n^{(\text{NP})}|^2$ term is also included in Eq. (4). In what follows, we will be neglecting QCD running of the local operators generated by the NP interaction (*i.e.* $\bar{\mathcal{C}}_1 = 0$ and $\bar{\mathcal{C}}_2 = 1$).

Models with extra vector-like quarks: Consider a model of the above type which extends the SM by including new 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 denote flavor and color respectively). Then the flavor-changing couplings are described by

$$\mathcal{L}_{xQ} = -\frac{g}{2\cos\theta_W} J_\mu^{(\text{NC})} Z^\mu + \text{h.c.} \quad , \quad (14)$$

$$J_\mu^{(\text{NC})} = \mathbf{U}_{ab}^{(u)} \bar{U}_{a,i} \Gamma_\mu U_{b,i} + \mathbf{U}_{ab}^{(d)} \bar{D}_{a,i} \Gamma_\mu D_{b,i} \quad ,$$

with flavor-changing couplings in both up and down sectors. The lifetime difference for this model can be obtained from Eq. (12) by substituting $\lambda_{sd} = \mathbf{U}_{cu}^{(u)} \mathbf{U}_{sd}^{(d)\dagger}$, $\lambda_{ds} = \lambda_{dd} = \lambda_{ss} = 0$ and $\Lambda = \sqrt{2/G_F}$. This model is well-constrained from measurements of the mass differences in $K\bar{K}$ and $D\bar{D}$ mixing. For $\mathbf{U}_{cu}^{(u)} \sim 10^{-3}$ and $\mathbf{U}_{sd}^{(d)} \sim 10^{-4}$ we get $y \sim 10^{-8}$, of the same order of magnitude as $y_{\text{SM}}^{(\text{LO})}$ 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 in the flavor $SU(3)$ limit is SUSY without R-parity [11]. In this

model, there are flavor-changing interactions of sleptons that can be obtained from the lagrangian

$$\mathcal{L}_{\mathcal{R}} = \lambda'_{ijk} L_i Q_j D_k^c \quad , \quad (15)$$

as well as the interactions mediated by squarks discussed below. The slepton-mediated interaction is not suppressed in the flavor $SU(3)$ limit and leads to

$$y_{\mathcal{R}} = \frac{C'\tilde{\lambda}}{M_\ell^2} \left[(C_2 - 2C_1) \langle Q' \rangle + (C_1 - 2C_2) \langle \tilde{Q}' \rangle \right], \quad (16)$$

where $C' = -G_F m_c^2 / (6\sqrt{2}\pi M_D \Gamma_D)$, M_ℓ is a slepton mass, $\tilde{\lambda}$ is given by Eq. (13) with $\lambda_{sd} = \lambda'_{i12} \lambda'_{i21} \leq 1 \times 10^{-9}$, $\lambda_{ss} = \lambda'_{i11} \lambda'_{i21} \leq 5 \times 10^{-5}$, $\lambda_{dd} = \lambda'_{i21} \lambda'_{i22} \leq 5 \times 10^{-5}$, $\lambda_{ds} = \lambda'_{i11} \lambda'_{i22} \leq 5 \times 10^{-2}$ [11], and $\langle Q' \rangle$ is

$$\langle Q' \rangle = \langle \bar{D}^0 | \bar{u}_i \gamma_\mu P_L c_i \bar{u}_j \gamma^\mu P_R c_j | D^0 \rangle \quad . \quad (17)$$

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 $M_\ell = 100$ GeV, we arrive at $y_{\mathcal{R}} \simeq -3.7\%$. 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 ($W^{(\text{R})}$) bosons [12]. The relevant effective lagrangian is

$$\mathcal{L}_{\text{LR}} = -\frac{g_{\text{R}}}{\sqrt{2}} \mathbf{V}_{ab}^{(\text{R})} \bar{u}_{a,i} \gamma^\mu P_R d_{b,i} W_\mu^{(\text{R})} + \text{h.c.} \quad , \quad (18)$$

where $\mathbf{V}_{ik}^{(\text{R})}$ are the coefficients of the right-handed CKM matrix. This leads to a local $\Delta C = -1$ hamiltonian as in Eq. (6) with $\bar{\Gamma}_1 = \bar{\Gamma}_2 = \gamma_\mu P_R$. Since current experimental limits allow $W^{(\text{R})}$ masses as low as a TeV [3], a sizable contribution to y is quite possible. Using Eq. (7), we obtain

$$y_{\text{LR}} = -C_{\text{LR}} \mathbf{V}_{cs}^{(\text{R})} \mathbf{V}_{us}^{(\text{R})*} \left[\mathcal{C}_1 \langle Q' \rangle + \mathcal{C}_2 \langle \tilde{Q}' \rangle \right] \quad , \quad (19)$$

where $C_{LR} \equiv \lambda G_F^{(R)} G_F m_c^2 x_s / (\pi M_D \Gamma_D)$, $G_F^{(R)} / \sqrt{2} \equiv g_R^2 / 8M_{WR}^2$, $C_{1,2}$ 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' ($\mathbf{V}^{(L)} = \mathbf{V}^{(R)}$) gives $y_{LR} = -4.8 \cdot 10^{-6}$ with $M_{WR} = 1.6$ TeV and (ii) 'Non-manifest LR' ($\mathbf{V}_{ij}^{(R)} \sim 1$) gives $y_{LR} = -8.8 \cdot 10^{-5}$ with $M_{WR} = 0.8$ 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 provides new tree-level contributions mediated by charged Higgs bosons and leads to the local four fermion interaction [13]

$$\mathcal{H}_{\text{ChH}}^{\Delta C=-1} = -\frac{\sqrt{2}G_F}{M_H^2} \bar{u}_i \bar{\Gamma}_1 q'_i \bar{q}_j \bar{\Gamma}_2 c_j \quad , \quad (20)$$

where the vertices $\bar{\Gamma}_{1,2}$ are

$$\begin{aligned} \bar{\Gamma}_1 &= m_{q'} \cot \beta V_{uq'} P_R - m_u \tan \beta V_{uq'} P_L \quad , \\ \bar{\Gamma}_2 &= m_q \cot \beta V_{cq}^* P_L - m_c \tan \beta V_{cq}^* P_R \quad . \end{aligned} \quad (21)$$

There are four possible contributions involving the various terms in $\bar{\Gamma}_{1,2}$. However, three of these, including the potentially large $\tan^2 \beta$ term, vanish for assorted reasons (*e.g.* flavor cancellation, zero matrix element). This leaves

$$y_{\text{ChH}} = \frac{\lambda^2 G_F^2 m_c^4 x_s^{3/2}}{\pi M_D \Gamma_D M_H^2} \cot^2 \beta [C_1 + C_2] \langle \mathcal{Q} \rangle \quad , \quad (22)$$

where $\langle \mathcal{Q} \rangle$ is as in Eq. (11). Assuming values $M_H = 85$ GeV and $\cot \beta = 0.05$, consistent with constraints obtained from the observation of $B \rightarrow \tau \nu_\tau$ [3], we obtain $y_{\text{ChH}} \simeq 2 \cdot 10^{-10}$.

SUSY without R-parity (squark exchange): The baryon-number violating squark exchanges arise from the lagrangian [11]

$$\mathcal{L}_{\mathcal{R}} = \lambda''_{ijk} U_i^c D_j^c D_k^c \quad . \quad (23)$$

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

$$y'_{\mathcal{R}} = -x_s C'' \frac{\lambda''_{22k} \lambda''_{12k}}{M_{sq}^2} [C_2 \langle Q' \rangle + C_1 N_c \langle \tilde{Q}' \rangle] \quad , \quad (24)$$

where $C'' = G_F \lambda / (2m_D \pi \Gamma)$, M_{sq} is a squark mass and the matrix elements $\langle Q' \rangle$, $\langle \tilde{Q}' \rangle$ are given earlier. Using factorization for the matrix elements, $\lambda''_{22k} \lambda''_{12k} \sim 3 \cdot 10^{-4}$ [14], and taking $M_{sq} = 100$ GeV, we arrive at the result $y'_{\mathcal{R}} \simeq 6.4 \cdot 10^{-6}$.

In conclusion, we have explored how NP contributions can influence the lifetime difference y 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 $D^0 \bar{D}^0$ 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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