Impact of colored scalars on $D^0-\bar{D}^0$ mixing in diquark models

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Inspired by the recent observation of the $D^0-\bar{D}^0$ mixing, we explore the effects of colored scalars on the $\Delta C = 2$ process in diquark models. As an illustration, we investigate the diquarks with the quantum numbers of $(6, \ 1, 1/3)$ and $(6, \ 1, 4/3)$ under $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge symmetries, which contribute to the process at one-loop and tree levels, respectively. We show that $\Delta m_D$ gives the strongest constraint on the free parameters. In addition, we find that the small couplings can be naturally interpreted by the suppressed flavor mixings if the diquark of $(6, \ 1, 4/3)$ only couples to the third generation.

A B S T R A C T

Unlike $K$ and $B$ systems, the short-distance (SD) contributions to charmed-meson flavor changing neutral current (FCNC) processes, such as the $D^0-\bar{D}^0$ mixing [1] and the decays of $c \to u\ell^+\ell^-$ and $D \to \ell^+\ell^-$ [2], are highly suppressed due to the stronger Glashow–Iliopoulos–Maiani (GIM) mechanism [3] and weaker heavy quark mass enhancements in the loops. In addition, it is often claimed that the long-distance (LD) effect for the $D^0-\bar{D}^0$ mixing should be the prevailing contribution in the SM. Nevertheless, because the nonperturbative hadronic effects are hard to control, the issue of the LD dominance is still inconclusive [4–7]. Recently, besides the progress of observing the $B_s$ oscillation [8–12], the evidence for the $D^0-\bar{D}^0$ mixing has also been exposed with the world averages [13]

\[ x_D = \frac{\Delta m_D}{T_D} = (0.811 \pm 0.334)\% , \]

\[ y_D = \frac{\Delta f_D}{2T_D} = (0.309 \pm 0.281)\% , \]

\[ y_{DCP} = (1.072 \pm 0.257)\% , \]

where $x_D$ ($y_D$) denotes the mass (lifetime) difference parameter and $y_{DCP}$ is the mixing parameter including the CP violating information. That is, if no CP violation is found in the $D$-meson oscillation, we have $y_{DCP} = y_D$. Due to these data, lots of studies on the physics beyond Standard Model (SM) have been done [14–20]. In particular, the possible extensions of the SM for the $D^0-\bar{D}^0$ mixing have been investigated by the authors in Ref. [21] in great detail.

In this note, we explore the issue in scalar diquark models which were not included in the previous discussions [21]. In the literature, the motivation to study the light colored scalar could be traced to the solution for the strong CP problem [22], where to avoid the domain-wall problem on the spontaneous CP violating mechanism, the models were constructed in the framework of grand unified theories (GUTs), e.g. $SU(5)$ gauge symmetry. The associated new source of CP violation on $K$ and $B$ systems was also studied in Refs. [23,24]. In addition, since the scalar sector in the SM has not been tested experimentally, it is plausible to assume the existence of other possible scalars in the gauge symmetry of $SU(3)_C \times SU(2)_L \times U(1)_Y$. Accordingly, the general scalar representations could be $(1, \ 2, 1/2), \ (8, \ 2, 1/2), \ (6, \ 3, 1/3), \ (6, \ 3, 4/3, 1/3, \ 2/3, \ 3, \ 3, \ -1/3, \ 3, \ 1, 2/3, -1/3, -4/3, \ -3/2, -3)$. The first (second) argument in the brackets denotes the CP representation in color (weak isospin) space and the number in the subscript corresponds to the hypercharge of $U(1)_Y$ [25]. Besides the SM Higgs doublet, it has been shown in Ref. [25] that when the hypothesis of minimal flavor violation (MFV) is imposed, only the representation $(8, \ 2, 1/2)$ could avoid FCNCs at tree level. As a result, due to the suppression of Cabibbo–Kobayashi–Maskawa (CKM) matrix elements and the masses of light quarks, the loop induced $D^0-\bar{D}^0$ oscillation in the color octet model is also negligible. Therefore, in the following analysis, we will concentrate on the situation of color triplet and sextet. In terms of involved Feynman diagrams, we find that $\Delta m_D$ can be produced by the diquark models through both box and tree diagrams. The various possible scalar diquarks are presented in Table 1, where the second column in the table denotes the representations of the diquarks under $SU(3)_C \times SU(2)_L \times U(1)_Y$, the third column gives the interactions of quarks and diquarks, the fourth column displays the relation of couplings in flavors and the last column shows the type of the effect that generates $\Delta m_D$. From the table, we see that only Model (7) can lead to the...
ΔC = 2 interaction at tree level. Due to the antisymmetric property in flavor indices, Model (6) cannot contribute to the ΔC = 2 process at tree level. It is worth mentioning that the colored sextet scalar also exists in a class of partial unification theories based on SU(2)L × SU(2)R × SU(4)C [26].

Since our purpose is to show the influence of diquarks on the D0–D̄0 oscillation, we are not planning to calculate the contributions of each model shown in Table 1. For comparison, we use Models (4) and (7), which have the same color structure, to illustrate the diquark effects. It is expected that the contributions of other models should be similar in order of magnitude.

1. Diquark (6, 1, 1/3)

We first write the interaction of quarks and the diquark of (6, 1, 1/3) as

\[ \mathcal{L}_{H_D} = f_{ij}^D d_{i\alpha}^C P_t u_{j\beta}^C H_D^{\alpha\beta} + \text{h.c.}, \]  

(2)

where \( f_{ij}^D \) denote the couplings of diquark and various flavors, \( C = i\gamma^5 \gamma^2 \) is the charge-conjugation operator, \( P_{t,\bar{t}}(R) = (1 \mp \gamma_5)/2 \) is the chiral projection operator and \( H_D^{\alpha\beta} \) is a weak gauge-singlet and colored sextet scalar with \( \alpha \) and \( \beta \) being the color indices.

After Fierz transformation, since the structure of four-fermion interactions becomes \( \delta \Gamma d \Gamma \Gamma \Gamma e \) completely lifetime and mass differences in the D0–D̄0 mixing cannot be induced at tree level. However, they can be produced at one-loop where the box diagrams are sketched in Fig. 1.

To formulate the ΔC = 2 effective Hamiltonian, we set the flavor indices in Fig. 1 to be \( k = \ell = c \) and \( j = n = u \). Hence, by including Wick contractions and neglecting the external momenta and the internal masses of light quarks, the effective Hamiltonian for Fig. 1 is written as

\[ -i H_{\Delta C=2} = \frac{1}{2} C_D^2 (\delta^\alpha_\rho \delta^\rho_\beta + \delta^\alpha_\beta \delta^\rho_\rho) \times \int \frac{d^4q}{(2\pi)^4} \frac{q^\mu q^\nu}{(q^2)^2(q^2 - m_H^2)^2} \bar{u}_\beta \gamma_\mu P_R c_\rho \bar{u}_\rho \gamma_\mu P_R c_\sigma, \]

(3)

with \( C_D = \sum f_{ij}^D f_{ju} \). Here, we have used the propagator for the sextet scalar as

\[ \langle TH_{\alpha\beta} H^{\nu\sigma}\rangle = \int \frac{d^4q}{(2\pi)^4} e^{-iq\cdot x} (\delta^\nu_\beta \delta^\sigma_\alpha + \delta^\nu_\alpha \delta^\sigma_\beta). \]

(4)

\[ \text{Fig. 1: Diquark box diagrams for the D0–D̄0 mixing.} \]

\[ \text{Fig. 2: Diquark-mediated flavor diagram for D+ → π+ (π0, \phi).} \]

With the loop integral

\[ \int \frac{d^4q}{(2\pi)^4} \frac{q^\mu q^\nu}{(q^2)^2(q^2 - m_H^2)^2} = -i \frac{m_H^{\mu\nu}}{4(4\pi)^2 m_H^2}. \]

(5)

Eq. (3) can be expressed as

\[ \mathcal{H}_{\Delta C=2} = \frac{C_D^2}{128\pi^2 m_H^2} \left[ 5 \bar{u}_\beta \gamma_\mu P_R c_\rho \bar{u}_\rho \gamma_\mu P_R c_\beta + \bar{u}_\beta \gamma_\mu P_R c_\beta \gamma_\mu P_R c_\beta \right]. \]

Using the transition matrix element given by

\[ \langle \bar{D}^0 | \bar{u}_\beta \gamma_\mu c_\rho \bar{u}_\rho \gamma_\mu c_\beta | D^0 \rangle = \frac{2}{3} f_D^2 m_D^2 B_D, \]

(6)

the

\[ x_D = \frac{\Delta m_D}{f_D} = Z_D \left( \frac{C_D}{m_H} \right)^2, \]

(7)

with \( \Delta m_D = 2|M_{12}| = 2|\langle \bar{D}^0 | \mathcal{H}_{\Delta C=2} | D^0 \rangle| \). Taking \( f_D = 4.10 \times 10^{-3} \), \( m_D = 1.86 \) GeV, \( f_{D^0} = 0.222 \) GeV [28] and \( B_D = 0.82 \) [29], in order to fit the current experimental data, the unknown parameters should satisfy

\[ \left( \frac{C_D}{m_H} \right)^2 = x_D \frac{Z_D}{\Delta m_D} = 4.7 \times 10^{-7} x_D \text{ GeV}^{-2}. \]

(8)

With \( m_H \sim 1 \) TeV and \( x_D \sim 8 \times 10^{-3} \), the free parameter \( |C_D| \) is about 0.06.

Besides a serious constraint on the free parameters from \( \Delta m_D \), for comparison, we consider other possible limits from rare nonleptonic D decays, such as \( D^+ → \pi^+ \pi^0 \) and \( D^+ → \pi^0 \phi \) decays, which are Cabibbo-suppressed and Cabibbo and color-suppressed processes in the SM, respectively. Their current measurements are [28]

\[ B(D^+ → \pi^+ \pi^0) = (1.24 \pm 0.07) \times 10^{-3}, \]

\[ B(D^+ → \pi^0 \phi) = (6.2 \pm 0.7) \times 10^{-3}. \]

(9)

According to the interactions in Eq. (2), flavor diagrams for D decays are given in Fig. 2 and the corresponding interactions for \( e → udd(\bar{ss}) \) are found to be

\[ \mathcal{H}_{e \rightarrow u} = -\frac{f_{q\bar{q}^\prime} f_{uu}}{2m_{tH}} (O_1^q + O_2^q), \]

(10)

where \( O_1^q = \bar{u}_\mu \gamma_\mu P_R c_\rho \gamma_\mu P_R c_\sigma \) and \( O_2^q = \bar{u}_\mu \gamma_\mu P_R c_\rho \bar{u}_\sigma \gamma_\mu P_R c_\omega \).
with \( q = d \) and \( s \). Based on the decay constants and transition form factor, defined by

\[
\langle 0 | \bar{q}^c \gamma^\mu \gamma_5 q | p (p) \rangle = i f_p p^\mu, \quad \langle 0 | \bar{q}^c \gamma^\mu s | q (p) \rangle = i m_\phi f_\phi \varepsilon_\phi^\mu (p),
\]

\[
\langle \pi (p_2) | \bar{u} \gamma_5 u | D (p_1) \rangle = f_+ (k^2) \left( p_\mu - \frac{p \cdot k}{k^2} k_\mu \right) + \frac{p \cdot k}{k^2} f_0 (k^2) k_\mu,
\]

(11)

with \( p = p_1 + p_2 \) and \( k = p_1 - p_2 \), the decay amplitudes by the naive factorization approach for \( D^+ \to \pi^+ (\pi^0, \phi) \) are given by

\[
A(\pi^+ \pi^0) = -i \frac{f_{d c} f_{d u}}{8 \sqrt{2} m_H^2} \left( 1 + \frac{1}{N_c} \right) f_\pi f_0 (0) m_D^2,
\]

\[
A(\pi^+ \phi) = -i \frac{f_{d c} f_{d u}}{8 \sqrt{2} m_H^2} \left( 1 + \frac{1}{N_c} \right) f_\phi f_+(0) p_1 \cdot \varepsilon_\phi.
\]

(12)

Here, we have taken the approximation of \( m_\ell^2 \approx 0 \) and set \( N_c = 3 \).

As a result, the branching ratios (BRs) are known as

\[
B(D^+ \to \pi^+ \pi^0) = \frac{\tau_D m_D^2}{211 \pi} \left( 1 + \frac{1}{N_c} \right)^2 \frac{f_{d c} f_{d u}}{m_H^2},
\]

\[
B(D^+ \to \pi^+ \phi) = \frac{\tau_D m_D^2}{211 \pi} \left( 1 + \frac{1}{N_c} \right)^2 \frac{f_\phi f_+(0)}{m_H^2} \frac{f_{d c} f_{d u}}{m_H^2}.
\]

(13)

For simplicity, we use the central values of the data to obtain the upper limits of the parameters. With \( f_\pi = 0.13 \) GeV, \( f_\phi = 0.237 \) GeV and \( f_+(0) = f_0 (0) = 0.624 \) [30], we get

\[
\left| \frac{f_{d c} f_{d u}}{m_H^2} \right|^2 < 1.7 \times 10^{-10} \text{ GeV}^{-4},
\]

\[
\left| \frac{f_{d c} f_{d u}}{m_H^2} \right|^2 < 1.5 \times 10^{-9} \text{ GeV}^{-4}.
\]

(14)

By comparing with Eq. (8), we see clearly that unless there exist strong cancelations among the free parameters, the constraints from the \( D^0 - \bar{D}^0 \) mixing is much stronger than those from \( D \) decays.

By examining Fig. 1, it is easy to find that down-type quarks involve in the internal loop. In other words, the \( K^0 - \bar{K}^0 \) mixing, denoted by \( \Delta m_{K} \), might give a strict constraint on the parameters. To understand the influence of \( \Delta m_{K} \), by the similar calculations on \( \Delta m_{D} \), the formula is given by

\[
\Delta m_K = 2 \text{Re} [K^0 | H_{\Delta m = 2} | K^0] = \frac{19}{768 \alpha^2} \frac{f_{d c} f_{d u} B_K}{m_H^2} \left( \frac{\text{Re} C_K}{m_H^2} \right)^2,
\]

(15)

with \( C_K = \sum_{j=u,c,t} f_{d j} f_{d j}^* \), where we have set \( m_t \ll m_H \). Clearly, although some parameters such as \( f_{u c} f_{d c}^* \) and \( f_{u t} f_{d t}^* \) appear in both \( \Delta m_{D} \) and \( \Delta m_{K} \), in general the \( C_K \) and \( C_D \) are different parameters. Adopting this viewpoint, the constraint from \( \Delta m_{K} \) might not have an influence on the constraint from \( \Delta m_{D} \). Nevertheless, in some special case, \( C_K \) and \( C_D \) are strongly correlated. Therefore, it is interesting to survey both constraints in a little bit of detail. Hence, according to Eqs. (7) and (15), we get

\[
\Delta m_K \approx \frac{f_{d c} f_{d u} B_K C_K^2}{m_H^2} \approx 0.14 \left| \frac{C_K}{C_D} \right|^2 m_H^2.
\]

(16)

With the data of \( \Delta m_{D} = x_D \Gamma_D \approx 0.008 \Gamma_D \approx 1.3 \times 10^{-14} \text{ GeV} \) and \( \Delta m_{K} = 3.483 \times 10^{-15} \text{ GeV} \) [28], the constrained relation from \( D \) and \( K \) systems is

\[
\left| \frac{C_K}{C_D} \right| \approx 1.4.
\]

(17)

Clearly, the result implies that when \( |C_K| \) and \( |C_D| \) are not regarded as independent parameters, both \( \Delta m_K \) and \( \Delta m_D \) will give similar constraints on the free parameters.

2. Diquark (6, 1, 4/3)

Next, we consider the contributions of the diquark (6, 1, 4/3), where the couplings to up-type quarks are written by

\[
\mathcal{L}_{H_6} = f_{ij} u_{ia}^T C_P c_{ij} u_{ja} H_6^{a | b} + \text{h.c.}
\]

(18)

Here, we adopt the same notations as those used for (6, 1, 1/3). However, one can find that the behavior of couplings \( f_{ij} \) is symmetric in flavors. To illustrate the result, the derivation is given as follows:

\[
f_{ij} u_{ia}^T C_P c_{ij} u_{ja} H_6^{a | b} = -f_{ji} u_{ia}^T C_P c_{ji} u_{ja} H_6^{a | b} = f_{ji} u_{ia}^T C_P c_{ji} u_{ja} H_6^{a | b},
\]

(19)

where we have used \( H_6^{a | b} = H_6^{a | b} \) in the last equality. Clearly, we get \( f_{ij} = f_{ji} \). The flavor diagrams for \( D \) decays are displayed in Fig. 3. To produce the \( \Delta C = 2 \) process, the flavor indices in Fig. 3 are chosen to be \( i = j = c \) and \( m = n = u \). With Wick contractions, the associated four-fermion interactions are obtained as

\[
\mathcal{H}_{\Delta C = 2} = -\frac{f_{c c} f_{u a}}{4 m_H^2} [\bar{u}^c \gamma_\mu P_R u_j \gamma^{i | b}] \bar{u}_R \gamma_\mu P_R c_j \bar{u}_a \gamma^{i | a} P_R c_R b.
\]

(20)

Using the transition matrix element of Eq. (6), we find

\[
\langle \bar{D}^0 | H_{\Delta C = 2} | D^0 \rangle = -\frac{7}{24} f_D^2 m_D^2 B_D \frac{f_{c c} f_{u a}}{m_H^2}.
\]

(21)

Consequently, the mixing parameter of the \( D^0 - \bar{D}^0 \) mixing is

\[
x_D = \frac{\text{Re} (f_{c c} f_{u a})^*}{m_H^2} \approx \frac{7}{24} f_D^2 m_D^2 B_D.
\]

(22)

From the values taken previously, the bound on the free parameters is found as

\[
|\text{Re} (f_{c c} f_{u a})^*| \frac{m_H^2}{m_D^2} = 7.2 \times 10^{-11} x_D \text{ GeV}^{-2}.
\]

(23)

If we adopt \( m_H \approx 1 \text{ TeV} \) and \( x_D \approx 8 \times 10^{-3} \), the constraint on the parameter is \( |\text{Re} (f_{c c} f_{u a})^*| \approx 5.76 \times 10^{-7} \). In particular, for \( f_{c c} \sim f_{u a} \) one gets \( |f_{c c}| \sim |f_{u a}| \approx 7.6 \times 10^{-3} \).

At the first sight, it seems that the small couplings are fine-tuned. However, the smallness in fact could be related to the suppressed flavor mixings. To demonstrate the conjecture, we propose that before the electroweak symmetry breaking, the scalar diquark only couples to one flavor, such as the top quark. Accordingly, the interaction is set to be

\[
\mathcal{L}_{H_6} = f_{H_6} H_6^{a | b} + \text{h.c.}
\]

(24)
After the symmetry breaking, to diagonalize the up-quark mass matrix, we introduce unitary matrices $V_L$ and $V_R$ so that the physical and weak eigenstates are related by $u_{RL} = V_{RL} u_{WRL}$. Using the relation, the couplings in Eq. (18) can be related to $f_i$ of Eq. (24) and $V_R$ by

$$f_{ij} = f_i (V_{Rij} V_{Rji})^*. \quad (25)$$

If we take $V_{R23} \sim \lambda^2$ and $V_{R13} \sim \lambda^3$ with $\lambda \sim 0.22$, then $f_{cc} f_{uu}^* = f_{tc}^2 V_{R3}^2 V_{R13} \sim f_{tt}^2 f_{uu}^2 = 2.7 f_{tt}^2 \times 10^{-7}$. By choosing suitable value for $f_t$, the proposed scenario in the analysis can naturally fit the result $|\text{Re}(f_{cc} f_{uu}^*)| \sim 5.76 \times 10^{-7}$ from the $\Delta D^0$ mixing. Intriguingly, the proposed model is similar to the case studied by the authors in Ref. [26] which is based on the $SU(2)_L \times SU(2)_R \times SU(4)_C$ symmetry. Moreover, the detailed studies on the production of colored sextet scalar at Colliders could be referred to Refs. [26,27].

In summary, we have investigated the contributions of the scalar diquarks to the $\Delta D^0$ oscillation. We have shown that the scalar diquark of (6, 1, 1/3) can generate the $\Delta C = 2$ process through box diagrams. By comparing with Cabibbo- and color-suppressed $D$ decays, the most strict limit on the free parameters arises from $\Delta m_D$. We also show that the constraints of $\Delta m_D$ and $\Delta m_K$ are comparable. For the diquark of (6, 1, 4/3), we have found that the $\Delta C = 2$ process can be induced at tree level. Although $\Delta m_D$ gives a very strong constraint on the free parameters, for the model with the diquark only coupling to the top-quark, the resultant small couplings can be ascribed to the small elements of the flavor mixing.

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