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CP violation in $D^0 \to (K^-K^+, \pi^-\pi^+)$ from diquarks

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Abstract

The explanation of the large CP asymmetries in $D^0 \to (\pi^+\pi^-, K^+K^-)$ decays observed by the LHCb collaboration is likely to call for new physics beyond the CKM paradigm. We explore new contributions caused by the color-sextet scalar diquark, and demonstrate that the diquark with the mass of order 1 TeV and nominal couplings with quarks can generate the CP asymmetries at the percent level. Using the experimental data on branching ratios and CP asymmetries of $D^0 \to (\pi^+\pi^-, K^+K^-)$, we derive the constraints on the diquark mass and couplings, which can be further examined in hadron colliders in the dijet final states.

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It is generally anticipated that both direct and indirect CP asymmetries (CPAs) in the charm sector are quite small in the standard model (SM). Any observation of the large CPA in D^0 decays will presumably imply that the underlying theory is out of the scope of the SM.

Recently based on the 0.62 fb^{-1} of data collected in 2011, the LHCb collaboration [1] has measured the difference between the time-integrated CP asymmetries in the decays $D^0 \to K^+K^-$ and $D^0 \to \pi^+\pi^-$, $\Delta A_{CP} \equiv A_{CP}(D^0 \to K^+K^-) - A_{CP}(D^0 \to \pi^+\pi^-)$, given by

$$\Delta A_{CP} = (-0.82 \pm 0.21(\text{stat.}) \pm 0.11(\text{sys.}))\% \,, \tag{1}$$

where the first uncertainty is statistical and the second is systematic. The quantity $A_{CP}(D^0 \to f)$ is defined as

$$A_{CP}(D^0 \to f) \equiv \frac{\Gamma(D^0 \to f) - \Gamma(\bar{D}^0 \to f)}{\Gamma(D^0 \to f) + \Gamma(\bar{D}^0 \to f)} \,, \tag{2}$$

with $f = K^+K^-, \pi^+\pi^-$. By contrast, results released by the CDF collaboration [2] based on 5.9 fb^{-1} of the integrated luminosity are somewhat less conclusive, given by

$$\begin{aligned} A_{CP}(D^0 \to \pi^+\pi^-) &= (+0.22 \pm 0.24 \pm 0.11)\% \,, \\ A_{CP}(D^0 \to K^+K^-) &= (-0.24 \pm 0.22 \pm 0.09)\% \,, \end{aligned} \tag{3}$$

while the previous world averages from Heavy Flavor Averaging Group [3] in 2010 are

$$\begin{aligned} A_{CP}(D^0 \to \pi^+\pi^-) &= (+0.22 \pm 0.37)\% \,, \\ A_{CP}(D^0 \to K^+K^-) &= (+0.16 \pm 0.23)\% \,. \end{aligned} \tag{4}$$

The new world average for ΔA_{CP} from Eqs. (1), (3) and (4) is found to be [4]

$$\Delta A_{CP} = -(0.645 \pm 0.180)\% \,, \tag{5}$$

which is about 3.6σ away from zero.

Contributions to $A_{CP}(D^0 \to f)$ contain both direct ($A_{CP}^{dir}(D^0 \to f)$) and indirect ($A_{CP}^{ind}(D^0 \to f)$) parts, and from the LHCb report [1] one has

$$\Delta A_{CP} \simeq \Delta A_{CP}^{dir} + (9.8 \pm 0.3)\% A_{CP}^{ind} \,, \tag{6}$$

where $\Delta A_{CP}^{dir} \equiv A_{CP}^{dir}(D^0 \to K^+K^-) - A_{CP}^{dir}(D^0 \to \pi^+\pi^-)$, and $A_{CP}^{ind} = A_{CP}^{ind}(D^0 \to f)$, which is universal for $f = K^+K^-$ and $\pi^+\pi^-$ and less than 0.3% due to the mixing parameters. Clearly, the LHCb data in Eq. (1) is dominated by the difference of the direct CP asymmetries, ΔA_{CP}^{dir}.

In order to have a nonzero direct CPA, two amplitudes A_1 and A_2 with both nontrivial weak phase difference θ_W and strong phase difference δ_S are called for, leading to

$$\begin{aligned} A_{CP}^{dir}(D^0 \to f) &= \frac{|A_f|^2 - |\bar{A}_f|^2}{|A_f|^2 + |\bar{A}_f|^2} \\ &= \frac{-2|A_1||A_2|\sin\theta_W \sin\delta_S}{|A_1|^2 + |A_2|^2 + 2|A_1||A_2|\cos\theta_W \cos\delta_S} \\ &\simeq -2r_f \sin\theta_W \sin\delta_S, \end{aligned} \tag{7}$$

with $r_f = |A_2|/|A_1|$. In the last step, a hierarchy of $r_f \ll 1$ has been adopted. The SM description of the direct CPA for $D^0 \to f$ arises from the interference between tree and penguin contributions, in which decay amplitudes take the generic expressions

$$A_f = V_{cq}^* V_{uq} T_{\rm SM} + V_{cb}^* V_{ub} P_{\rm SM}, \tag{8}$$

with $q = s$ for $f = K^+K^-$ and $q = d$ for $\pi^+\pi^-$. Besides the hierarchy in the CKM matrix elements $V_{cq}^* V_{uq} \gg V_{cb}^* V_{ub}$, penguin amplitudes are also suppressed by loop factors. Even in the limit $P_{SM} \sim T_{SM}$, the ratio of the decay amplitudes is still very small $r_f \sim 0.0007$, leading to a tiny CPA which is far below the central value in Eq. (5). As a result, if we took the data by the LHCb seriously, a solution to the large ΔA_{CP} would be to introduce some new CP violating mechanism beyond the CKM.

By neglecting the small SM penguin contributions, the decay amplitude of $D^0 \to f$ ($f = K^+K^-$ or $\pi^+\pi^-$) with new physics contributions could be parametrized as

$$A_f = V_{cq}^* V_{uq} \left[T'_{SM} \left(1 + \rho e^{i\theta_W} \right) + E'_{SM} e^{i\delta_S} \right] , \tag{9}$$

where θ_W is the new weak CP phase ranging from 0 to π, while ρ is associated with the new physics effect with an arbitrary sign, *i.e.*, sign(ρ) $= \pm 1$. In the above equation, T'_{SM} corresponds to the W emission diagram, while E'_{SM} is from the annihilation type of the W-exchange diagram. We note that since the final state interactions make dominant contributions to the W-exchange diagram, without losing generality, we regard that the short-distance (SD) effect of new physics on such topology could be ignored. Consequently, in this circumstance new physics plays an important role for the emission topology. From Eqs. (7) and (9), we find

$$A_{CP}^{dir}(D^0 \to f) = \frac{4T'_{SM} E'_{SM} \rho \sin\theta_W \sin\delta_S}{|a(\zeta)|^2 + |a(\zeta \to \zeta^*)|^2} \tag{10}$$

with $a(\zeta) = T'_{SM}\zeta + E'_{SM}e^{i\delta_S}$ and $\zeta = 1 + \rho\exp(i\theta_W)$. If all quantities in the SM are under control, ρ and θ_W are the only free parameters.

Recently, a number of theoretical studies [5–17] have been performed to understand the LHCb and CDF data. In this brief report, we would like to use the scalar diquarks as the sources of new physics. Although the introduction of scalar diquarks in the literature has its physical reason, such as solving the strong CP problem [18–20], here we only explore their consequences in D-meson processes. The combinations of two quarks give many possible types of diquarks, such as $(3,1,-1/3)$, $(6,1,1/3)$, $(3,3,-1/3)$, $(3,1,-4/3)$, and $(6,1,4/3)$ under the standard group of $SU(3)_C \times SU(2)_L \times U(1)_Y$, among which the $(3,1,-1/3)$ and $(6,1,1/3)$ diquarks are very interesting for avoiding the strong correlation in flavor couplings and the strict constraint from the tree induced D-$\bar{D}$ mixing [21]. In what follows we will concentrate on the color sextet $H_6 = (6,1,1/3)$ to illustrate its effect on the direct CPAs in $D^0 \to K^+K^-, \pi^+\pi^-$, but a similar study could be applied to the color triplet boson.

We first write the interaction between quarks and H_6 as

$$\mathcal{L}_{H_{\mathbf{6}}} = f_{ij} d^{cT}_{i\alpha} C P_L u^c_{j\beta} H^{\alpha\beta}_{\mathbf{6}} + h.c.\,, \tag{11}$$

where f_{ij} denotes the couplings between the diquark and quark flavors, $C = i\gamma^0\gamma^2$ is the charge-conjugation matrix, $P_{L(R)} = (1 \mp \gamma_5)/2$ is the chiral projection operator, and $H^{\alpha\beta}_{\mathbf{6}}$ is a weak gauge-singlet and colored sextet scalar with α and β being the color indices. In terms of the interactions in Eq. (11), flavor diagrams for D decays are given in Fig. 1. After integrating out the highly virtual diquarks, the corresponding interactions for $c \to u\bar{d}d(\bar{s}s)$ are derived as

$$\begin{aligned} \mathcal{H}_{c\to u} &= -\frac{f^*_{qc} f_{qu}}{16 m_H^2}\,(O^q_1 + O^q_2)\,, \\ O^q_1 &= (\bar{u}q)_{V+A}(\bar{q}c)_{V+A}\,, \\ O^q_2 &= (\bar{u}_\alpha q_\beta)_{V+A}(\bar{q}_\beta c_\alpha)_{V+A}\,, \end{aligned} \tag{12}$$

with $(\bar{q}''q')_{V+A} = \bar{q}''\gamma^\mu(1+\gamma_5)q'$ and $q = d, s$. Based on the decay constants and transition form factors, defined by

$$\begin{aligned} \langle 0|\bar{q}'\gamma^\mu\gamma_5 q|P(p)\rangle &= i f_P p^\mu\,, \\ \langle P(p_2)|\bar{q}\gamma_\mu c|D(p_1)\rangle &= f_+^{DP}(k^2)\Big\{Q_\mu - \frac{Q\cdot k}{k^2}k_\mu\Big\} \\ &\quad + \frac{Q\cdot k}{k^2} f_0^{DP}(k^2)\, k_\mu\,, \end{aligned} \tag{13}$$

respectively, with $Q = p_1 + p_2$ and $k = p_1 - p_2$, the diquark contribution to $D \to \bar{f}$ is given by

$$\rho e^{i\theta_W} = \frac{\sqrt{2}}{16 m_H^2 G_F} \frac{f_{qc}^* f_{qu} a_1'}{V_{cq}^* V_{uq} a_1}\,, \tag{14}$$

where $a_1 = c_1(\mu) + c_2(\mu)(1/N_c + \chi(\mu))$ and $a_1' = 1 + (1/N_c + \chi(\mu))$ are the effective Wilson coefficients [22] with $\chi(\mu)$ being the nonfactorizable contribution. In the large N_c limit, as the nonfactorizable effect could be simplified as $\chi = -1/N_c$ [23], the nonfactorizable part will be smeared by the operator O_2^q.

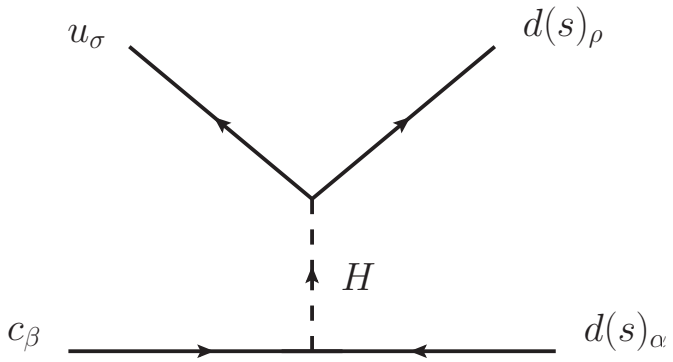

FIG. 1. Diquark-mediated flavor diagram for $D^0 \to (\pi^+\pi^-, K^+K^-)$

Before the numerical study, we discuss how to escape the stringent constraint from the $D - \bar{D}$ mixing. Using the results in Ref. [21], the D mixing parameter induced by box diagrams could be formulated by

$$x_D = \frac{\Delta m_D}{\Gamma_D} \sim \tau_D \frac{19 f_D^2 m_D}{1536\pi^2} \left| \frac{\sum_i f_{ic}^* f_{iu}}{m_H} \right|^2 . \tag{15}$$

Since the flavors (denoted by i) in the internal loops include d, s and b quarks, the constraint from x_D could be released if a cancellation occurs in $\sum_i f_{ic}^* f_{iu}$.

In our analysis, the input values of the SM are taken as [13, 22, 24]:

$$\begin{aligned}
&T'_{SM}(\pi\pi, KK) = (3.01, 4.0) \times 10^{-6}\ \text{GeV}\,,\\
&E'_{SM}(\pi\pi, KK) = (1.3, 1.6) \times 10^{-6}\ \text{GeV}\,,\\
&\delta_S(\pi\pi, KK) = (145, 108)^\circ\,, a_1 = 1.21,\\
&V_{us} = -V_{cd} = 0.2252\,, V_{cs} = 0.97345,\ \ V_{ud} = 0.97428\,,\\
&m_{\pi(K)} = 0.139(0.497)\ \text{GeV}\,, m_D = 1.863\ \text{GeV}\,,
\end{aligned} \tag{16}$$

where the resulting branching ratios (BRs) for $D^0 \to (\pi^-\pi^+, K^-K^+)$ are estimated as $(1.38, 3.96) \times 10^{-3}$, while the current data are $\mathcal{B}(D^0 \to \pi^-\pi^+) = (1.400 \pm 0.026) \times 10^{-3}$ and $\mathcal{B}(D^0 \to K^-K^+) = (3.96 \pm 0.08) \times 10^{-3}$ [24]. Since $f^*_{dc}f_{du}$ and $f^*_{sc}f_{su}$ are independent free parameters, to simplify our calculation, we adopt two benchmark schemes: (I) $f^*_{dc}f_{du} \approx f^*_{sc}f_{su}$ and (II) $f^*_{dc}f_{du} \approx -f^*_{sc}f_{su}$. Consequently, the involved parameters in the analysis are $\theta_W = arg(f^*_{sc}f_{su})$ and $|f^*_{sc}f_{su}|/m^2_H$.

An estimate of the scalar diquark contribution is given as follows. The current measurement of the dijet cross section from the hadron collider puts the limit of the scalar diquark mass, see Ref. [25, 26] for instance,

$$m_{H_6} > 1.9\text{TeV}, \tag{17}$$

where a normal diquark-quark coupling is used. When the diquark decay is taken into account, this value may get reduced. Assuming the mass of order 1 TeV and normal couplings for the diquark, we find from Eq. (14) that the NP contribution is at the percent level compared to the SM contribution. As a result, such a diquark is able to explain the large CPA data by the LHCb, while its effects to branching ratios will be up to a few percent.

Since the involved parameters in Cabibbo allowed processes, *e.g.*, $D \to \pi K$, are different from the singly Cabibbo suppressed decays, with the assumption of $\sum_i f^*_{ic}f_{iu} \to 0$, the BRs for $D^0 \to (\pi^+\pi^-, K^+K^-)$ are the potential constraint. Thus, we will take the data of the BRs with errors to constrain the free parameters. For Scheme I, we show the CPA difference defined in Eq. (6) as a function of $\xi \equiv \pm|f^*_{sc}f_{su}|/m^2_H$ and θ_W in Fig. 2. The solid lines correspond to the upper and lower bounds of the LHCb results. The dashed (blue) and dash-dotted (red) lines denote the experimental BRs for the decays $D^0 \to (\pi^+\pi^-, K^+K^-)$, respectively. From these curves, we obtain the allowed region of the parameters as

$$\begin{aligned} -2 \times 10^{-7} < \xi < -0.6 \times 10^{-7}, \\ 0.3 < \theta_W < 2.25. \end{aligned} \tag{18}$$

Similarly, the results in Scheme II are presented in Fig. 3. In this scheme, we find that the magnitude of ΔA^{dir}_{CP} cannot fit the data within 1σ errors. The reason is that the direct CPA of $D^0 \to \pi^+\pi^-$ in scheme II has the same sign with the one of $D^0 \to K^+K^-$. In order to understand the effects on each direct CPA A^{dir}_{CP} defined in Eq. (7), by using Scheme I, we display A^{dir}_{CP} for $D^0 \to \pi^+\pi^-$ and K^+K^- decays in Figs. 4a and 4b, respectively. The results, even the sign, are consistent with the current CDF data shown in Eq. (3).

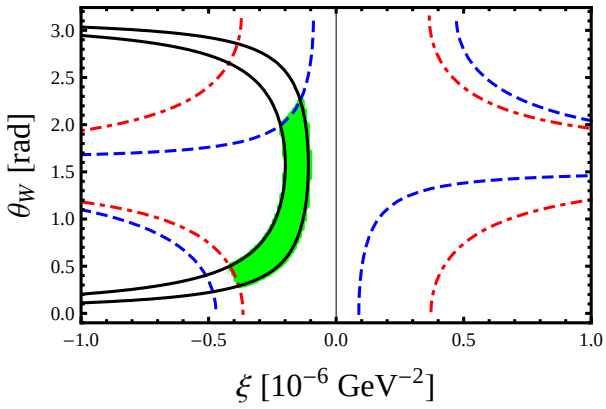

FIG. 2. ΔA_{CP}^{dir} (solid) in units of 10^{-2} as a function of $\xi \equiv \pm|f_{sc}^* f_{su}|/m_H^2$ and θ_W in the diquark model, where the dashed (blue) and dash-dotted lines denote the data with errors of BRs for the decays $D^0 \to (\pi^+\pi^-, K^+K^-)$ in units of 10^{-3}, respectively, while the shadowed region (green) is the combined constraint from BRs and direct CPAs.

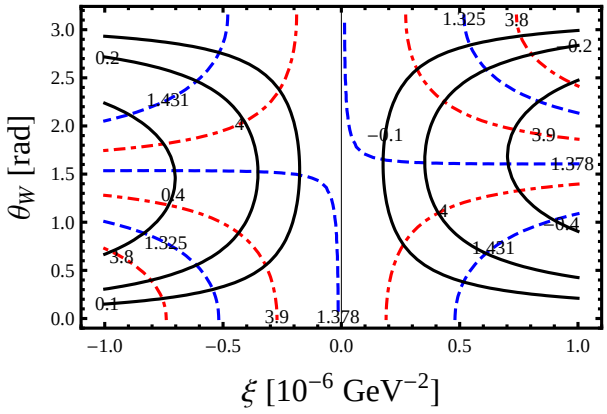

FIG. 3. Legend is the same as that in Fig. 2, but the magnitude of ΔA_C^{dir} is smaller than the LHCb results of 1σ.

In summary, it is clearly an exciting moment if any CPA in D-meson decays is observed at the percent level as the SM prediction is far below 1%. Motivated by the recent LHCb measurement on the time integrated CPA in $D^0 \to (\pi^+\pi^-, K^+K^-)$ decays, which appears to be inconsistent with the SM result, we have studied the contributions of the colored scalar boson and used the color sextet $(6, 1, 1/3)$ as the illustrator. In this diqaurk model, the serious constraint from the D mixing parameter induced by box diagrams could be avoided if a cancellation among different flavor couplings occurs. We have found that the the induced direct CPAs of $D^0 \to (\pi^+\pi^-, K^+K^0)$ decays can fit the recent LHCb results

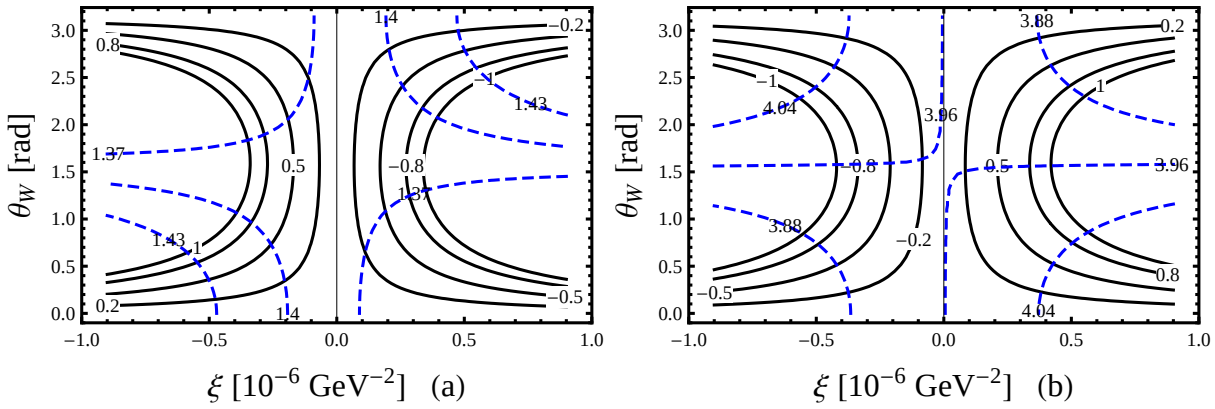

FIG. 4. Direct CPAs (solid lines) and BRs (dashed lines) for (a) $D^0 \to \pi^+\pi^-$ and (b) $D^0 \to K^+K^-$, where the units for the solid and dashed lines are 10^{-2} and 10^{-3}, respectively.

and are consistent with the CDF current measurement.

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