Pure annihilation type D→PP(V) decays in the perturbative QCD approach*

ZOU Zhi-Tian(邹芝田) LI Cheng(李承) LÜ Cai-Dian(吕才典)[1)]

Institute of High Energy Physics and Theoretical Physics Center for Science Facilities, Chinese Academy of Sciences, Beijing 100049, China

Abstract: The annihilation type diagrams are difficult to calculate in any kind of model or method. Encouraged by the successful calculation of pure annihilation type B decays in the perturbative QCD factorization approach, we calculate the pure annihilation type $D \to PP(V)$ decays in the perturbative QCD approach based on the k_T factorization. Although the expansion parameter $1/m_D$ is not very small, our leading order numerical results agree with the existing experimental data for most channels. We expect more accurate observation from experiments, which can help us learn about the dynamics of D meson weak decays.

Key words: D meson, perturbative QCD, annihilation

PACS: 13.20.Ft, 12.38.Bx, 14.40.Lb **DOI:** 10.1088/1674-1137/37/9/093101

1 Introduction

After decades of study, the D meson decays are still a hot topic on both the theoretical side and the experimental side, since they can provide useful information on flavor mixing, CP violation, strong interactions and even the new physics signal [1–3]. For example the recent observation of D^0-$\bar{D}^0$ mixing provides us a new platform to explore new physics via favor-changing neutral currents. To date, the CLEO-c and two B factories experiments have given many results about the D decays. The BES-III experiment is expected to give more results. The accurate observation can help us understand the QCD dynamics and the D meson weak decays. In recent years, many theoretical studies on the decays of the D meson have been done based on the diagrammatic approach [4], the final-state interaction effects [5, 6], a combination of factorization and pole model [7], the factorization-assisted topological diagrammatic approach [8], and the perturbative QCD (PQCD) approach [9].

Most of the theoretical studies show that the annihilation type diagrams in hadronic D decays play a very important role [4, 7–9]. For example in Ref. [4], the authors take the model-independent diagrammatic approach to study the two-body nonleptonic D decays, with all topological amplitudes extracted from the experimental data. Their analysis indicates that the $SU(3)$ breaking effect and the annihilation type contributions are important to explain the experimental data. The importance of the annihilation diagram contribution is also reflected from the large difference of D^0 and D^+ lifetime. However, these annihilation type diagrams are usually very difficult to calculate, since factorization may not work here. In Ref. [7], the authors use the pole model to give large annihilation diagram contributions. It is worth mentioning that the annihilation type diagrams can be perturbatively calculated without parametrization in the PQCD approach based on k_T factorization [10, 11]. For these pure annihilation type B decays, the predictions in the PQCD approach have been confirmed by experiments later [12–14].

The factorization that is proved in the $1/m_b$ expansion, can be applied to the corresponding D meson decays straightforwardly. However, the expansion is much poorer in D decays than that in B decays due to smaller D meson mass. Anyway, since there is no better method for the annihilation diagram calculation, the pure annihilation type decays $D^0 \to \bar{K}^0\phi$ were calculated in the PQCD approach [9], with a good agreement with the experimental result. In this work, we use the PQCD approach to analyze the 10 modes of pure annihilation type $D \to PP(V)$ decays. By keeping the intrinsic transverse momentum k_T of valence quarks, the end point singularity, which will spoil the perturbative calculation, can be regulated by Sudakov form factor and threshold resummation. Therefore, the PQCD approach can give converging results with predictive power.

In the standard model, two body hadronic D meson

Received 4 January 2013, Revised 24 January 2013

∗ Supported by National Science Foundation of China (11228512, 11235005, 11075168)

1) E-mail: lucd@ihep.ac.cn

weak decays are dominated by the contributions from tree operators, since the contributions from the penguin operators are suppressed both by the small elements of the Cabibbo-Kobayashi-Maskawa (CKM) matrix and by the relatively small b quark mass in the c-b-u penguin diagram. This is in contrast to the penguin amplitude in B decays, which can profit from a larger CKM element and a much larger t quark mass. Although the suppressed penguin diagram contributions may be the main source of the direct asymmetry [2, 3, 8, 15], we ignore the penguin contributions in this work due to the small effect on the branching fractions.

2 Formalism and perturbative calculation

For the pure annihilation type D→PP(V) decays, at the quark level, the dominant contributions are described by the effective Hamiltonian $H_{\rm eff}$

$$H_{\rm eff}=\frac{G_{\rm F}}{\sqrt{2}}V_{\rm uq}V^{*}_{\rm cq'}\left[C_1(\mu)O_1(\mu)+C_2(\mu)O_2(\mu)\right], \qquad (1)$$

where $V_{\rm cq'}$ and $V_{\rm uq}$ are the corresponding CKM matrix elements, with $\rm q^{(\prime)}=d$, s, and $C_{1,2}(\mu)$ are Wilson coefficients at the renormalization scale μ. $O_{1,2}(\mu)$ are the four quark operators from tree diagrams

$$O_1=(\bar{q}'_{\alpha}c_{\beta})_{V-A}(\bar{u}_{\beta}q_{\alpha})_{V-A},O_2=(\bar{q}'_{\alpha}c_{\alpha})_{V-A}(\bar{u}_{\beta}q_{\beta})_{V-A},$$

where α and β are the color indices, $(\bar{q}'_{\alpha}c_{\beta})_{V-A}=\bar{q}'_{\alpha}\gamma^{\mu}(1-\gamma^5)c_{\beta}$. Conventionally, the combination of Wilson coefficients can be defined as

$$a_1=C_2+C_1/3,a_2=C_1+C_2/3. \qquad (2)$$

In the hadronic matrix element calculation, the decay amplitude can be factorized into soft (Φ), hard (H), and harder (C) dynamics characterized by different scales [9, 16],

$$\begin{aligned}\mathcal{A}\sim&\int{\rm d}x_1{\rm d}x_2{\rm d}x_3b_1{\rm d}b_1b_2{\rm d}b_2b_3{\rm d}b_3\\&\times{\rm Tr}\big[C(t)\Phi_{\rm D}(x_1,b_1)\Phi_{\rm M_2}(x_2)\Phi_{\rm M_3}(x_3)\\&\times H(x_i,b_i,t)S_t(x_i){\rm e}^{-S(t)}\big], \qquad (3)\end{aligned}$$

where b_i is the conjugate space coordinate of quark's transverse momentum $k_{i\rm T}$, x_i is the momentum fractions of valence quarks, and t is the largest energy scale in the hard part function $H(x_i,b_i,t)$. $C(t)$ are the Wilson coefficients with resummation of the large QCD corrections of four quark operators. The large double logarithms $\ln^2x_i$ are summed by the threshold resummation to give a jet function $S_{\rm t}(x_i)$ which smears the end-point singularities on x_i [17]. The Sudakov form factor ${\rm e}^{-S(t)}$ is from resummation of double logarithms, which suppresses the soft dynamics effectively and the long distance contributions in the large b region [18, 19]. Thus it makes the perturbative calculation of the hard part H reliable. The meson wave functions Φ_i, are nonperturbative input parameters but universal for all decay modes.

The leading order Feynman diagrams of the considered decays are shown in Fig. 1. For D → PP decays, the amplitude from factorizable diagrams (a) and (b) in Fig. 1 is

$$\begin{aligned}\mathcal{A}_{\rm af}=&-8C_{\rm F}f_{\rm D}\pi m_{\rm D}^4\int_0^1{\rm d}x_2{\rm d}x_3\int_0^{1/\Lambda}b_2{\rm d}b_2b_3{\rm d}b_3\\&\times\Big\{\big[2\phi^{\rm P}_{\rm M_2}(x_2)r_{02}r_{03}(\phi^{\rm P}_{\rm M_3}(x_3)(x_3-2)-x_3\phi^{\rm T}_{\rm M_3}(x_3))\\&+\phi^{\rm A}_{\rm M_2}(x_2)\phi^{\rm A}_{\rm M_3}(x_3)(x_3-1)\big]h_{\rm af}(\alpha,\beta,b_2,b_3)E_{\rm af}(t_{\rm a})\\&+\big[2\phi^{\rm P}_{\rm M_3}(x_3)r_{02}r_{03}(\phi^{\rm T}_{\rm M_2}(x_2)(x_2-1)\\&+\phi^{\rm P}_{\rm M_2}(x_2)(x_2+1))+x_2\phi^{\rm A}_{\rm M_2}(x_2)\phi^{\rm A}_{\rm M_3}(x_3)\big]\\&\times h_{\rm af}(\alpha',\beta,b_3,b_2)E_{\rm af}(t_{\rm b})\Big\}, \qquad (4)\end{aligned}$$

where, $C_{\rm F}=4/3$ is the group factor of $SU(3)_{\rm c}$, and $r_{02(03)}=m_{02(03)}/m_{\rm D}$ with the chiral mass $m_{02(03)}$ of the pseudoscalar meson. The hard scale $t_{\rm e,f}$ and the functions $E_{\rm af}$ and $h_{\rm af}$ can be given by

$$\begin{aligned}t_{\rm a}=&\max\{\sqrt{(r_3^2+x_2(1-r_3^2))(1-r_2^2)(1-x_3)}m_{\rm D},\\&\sqrt{1-x_3(1-r_2^2)}m_{\rm D},1/b_2,1/b_3\},\\t_{\rm b}=&\max\{\sqrt{(1-r_2^2)(r_3^2+x_2(1-r_3^2))}m_{\rm D},\\&1/b_2,1/b_3\}, \qquad (5)\end{aligned}$$

$$E_{\rm af}(t)=\alpha_{\rm s}(t)\cdot\exp[-S_{\rm M_2}(t)-S_{\rm M_3}(t)], \qquad (6)$$

$$\begin{aligned}h_{\rm af}(\alpha,\beta,b_2,b_3)=&\left(\frac{{\rm i}\pi}{2}\right)^2H_0^{(1)}(\beta b_2)S_{\rm t}(x_3)\\&\times\Big[\theta(b_2-b_3)H_0^{(1)}(\alpha b_2)J_0(\alpha b_3)\\&+\theta(b_3-b_2)H_0^{(1)}(\alpha b_3)J_0(\alpha b_2)\Big], \qquad (7)\end{aligned}$$

with $r_{2(3)}=m_{\rm M_{2(3)}}/m_{\rm D}$, $\alpha^2=(1-x_3(1-r_2^2))m_{\rm D}^2$, $\beta^2=(r_3^2+x_2(1-r_3^2))(1-r_2^2)(1-x_3)m_{\rm D}^2$ and $\alpha'^2=(r_3^2+x_2(1-r_3^2))(1-r_2^2)m_{\rm D}^2$.

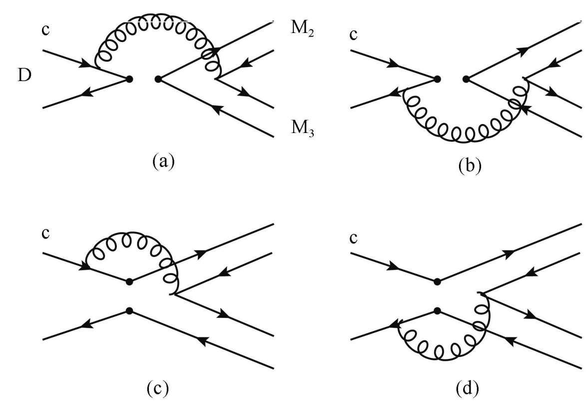

Fig. 1. The diagrams contributing to the pure annihilation type D→PP(V) decays in PQCD.

For the so-called non-factorizable diagrams (c) and (d) in Fig. 1, the decay amplitude is

$$\begin{aligned}\mathcal{M}_{\rm anf} =& 16\sqrt{\frac{2}{3}}C_{\rm F}\pi m_{\rm D}^4\int_0^1 {\rm d}x_1{\rm d}x_2{\rm d}x_3\int_0^{1/\Lambda} b_1{\rm d}b_1b_2{\rm d}b_2\\ &\times\phi_{\rm D}(x_1,b_1)\Big\{\big[\phi^{\rm A}_{\rm M_2}(x_2)\phi^{\rm A}_{\rm M_3}(x_3)(x_1+x_2)\\ &+r_{02}r_{03}\big(\phi^{\rm P}_{\rm M_2}(x_2)\big(\phi^{\rm P}_{\rm M_3}(x_3)(x_1+x_2-x_3+3)\\ &+\phi^{\rm T}_{\rm M_3}(x_3)(1-x_1+x_2-x_3)\big)\\ &+\phi^{\rm T}_{\rm M_2}(x_2)\ \big(\phi^{\rm P}_{\rm M_3}(x_3)(x_1+x_2+x_3-1)\\ &+\phi^{\rm T}_{\rm M_3}(x_3)(x_3-x_1-x_2+1)\big)\big)\big]\\ &\times h_{\rm anf1}(\alpha,\sqrt{|\beta_1^2|},b_1,b_2)E_{\rm anf}(t_{\rm c})\\ &+\big[\phi^{\rm A}_{\rm M_2}(x_2)\phi^{\rm A}_{\rm M_3}(x_3)(x_3-1)\\ &+r_{02}r_{03}\big(\phi^{\rm P}_{\rm M_2}(x_2)\big(\phi^{\rm P}_{\rm M_3}(x_3)(x_1-x_2+x_3-1)\\ &+\phi^{\rm T}_{\rm M_3}(x_3)(x_1-x_2-x_3+1)\big)\\ &+\phi^{\rm T}_{\rm M_2}(x_2)\big(\phi^{\rm P}_{\rm M_3}(x_3)(x_2+x_3-x_1-1)\\ &+\phi^{\rm T}_{\rm M_3}(x_3)(1-x_1+x_2-x_3)\big)\big)\big]\\ &\times h_{\rm anf2}(\alpha,\sqrt{|\beta_2^2|},b_1,b_2)E_{\rm anf}(t_{\rm d})\Big\},\end{aligned}\tag{8}$$

with

$$\begin{aligned}t_{\rm g} =& \max\{\sqrt{(r_3^2+x_2(1-r_3^2))(1-r_2^2)(1-x_3)}m_{\rm D},\\ &\sqrt{1-[(1-r_3^2)(1-x_2)-x_1][r_2^2+x_3(1-r_2^2)]}m_{\rm D},\\ &1/b_1,1/b_2\},\\ t_{\rm h} =& \max\{\sqrt{|(x_1-r_3^2-x_2(1-r_3^2))|(1-r_2^2)(1-x_3)}m_{\rm D},\\ &\sqrt{(r_3^2+x_2(1-r_3^2))(1-r_2^2)(1-x_3)}m_{\rm D},\\ &1/b_1,1/b_2\},\end{aligned}\tag{9}$$

$$E_{\rm anf}=\alpha_{\rm s}(t)\cdot\exp[-S_{\rm D}(t)-S_{\rm M_2}(t)-S_{\rm M_3}(t)]|_{b_2=b_3},\tag{10}$$

$$\begin{aligned}h_{\rm anfj} =& \frac{{\rm i}\pi}{2}\Big[\theta(b_1-b_2)H_0^{(1)}(\alpha b_1)J_0(\alpha b_2)\\ &+\theta(b_2-b_1)H_0^{(1)}(\alpha b_2)J_0(\alpha b_1)\Big]\\ &\times\begin{cases}\frac{{\rm i}\pi}{2}H_0^{(1)}\big(\sqrt{|\beta_j^2|}b_1\big), & \beta_j^2<0,\\ K_0\big(\sqrt{|\beta_j^2|}b_1\big), & \beta_j^2>0,\end{cases}\end{aligned}\tag{11}$$

where $j=1,2$, $\beta_1^2=1-[(1-r_3^2)(1-x_2)-x_1][r_2^2+x_3(1-r_2^2)]m_{\rm D}^2$, $\beta_2^2=(x_1-r_3^2-x_2(1-r_3^2))(1-r_2^2)(1-x_3)m_{\rm D}^2$, and $\alpha=\sqrt{(r_3^2+x_2(1-r_3^2))(1-r_2^2)(1-x_3)}m_{\rm D}$. The expressions of $S_{\rm D}(t)$, $S_{\rm M_2}(t)$, $S_{\rm M_3}(t)$ and $S_{\rm t}$ can be found in Refs. [17, 19, 20].

For those D→PV decays, the decay amplitudes are

$$\begin{aligned}\mathcal{A}^{\rm PV}_{\rm af} =& 8C_{\rm F}f_{\rm D}\pi m_{\rm D}^4\int_0^1{\rm d}x_2{\rm d}x_3\int_0^{1/\Lambda}b_2{\rm d}b_2b_3{\rm d}b_3\\ &\times\big\{\big[2\phi^{\rm P}_{\rm M_2}(x_2)r_{02}r_{\rm V}(\phi^{\rm s}_{\rm V}(x_3)(x_3-2)-x_3\phi^{\rm t}_{\rm V}(x_3))\\ &+\phi^{\rm A}_{\rm M_2}(x_2)\phi_{\rm V}(x_3)(r_{\rm V}^2-1)(x_3-1)\big]\\ &\times h_{\rm af}(\alpha,\beta,b_2,b_3)E_{\rm af}(t_{\rm a})\\ &-\big[-2\phi^{\rm s}_{\rm V}(x_3)r_{02}r_{\rm V}(\phi^{\rm T}_{\rm M_2}(x_2)(x_2-1)+(x_2+1)\\ &\times\phi^{\rm P}_{\rm M_2}(x_2))+(x_2+(1-2x_2)r_{\rm V}^2)\phi^{\rm A}_{\rm M_2}(x_2)\phi_{\rm V}(x_3)\big]\\ &\times h_{\rm af}(\alpha',\beta,b_3,b_2)E_{\rm af}(t_{\rm b})\big\},\end{aligned}\tag{12}$$

$$\begin{aligned}\mathcal{M}^{\rm PV}_{\rm anf} =& 16\sqrt{\frac{2}{3}}C_{\rm F}\pi m_{\rm D}^4\int_0^1{\rm d}x_1{\rm d}x_2{\rm d}x_3\int_0^{1/\Lambda}b_1{\rm d}b_1b_2{\rm d}b_2\\ &\times\phi_{\rm D}(x_1,b_1)\big\{\big[\phi^{\rm A}_{\rm M_2}(x_2)\phi_{\rm V}(x_3)\\ &(x_1+x_2+(-x_1-2x_2+x_3+1)r_{\rm V}^2))\\ &+r_{02}r_{\rm V}\big(\phi^{\rm T}_{\rm M_2}(x_2)(\phi^{\rm s}_{\rm V}(x_3)(1-x_1-x_2-x_3)\\ &+\phi^{\rm t}_{\rm V}(x_3)(x_1+x_2-x_3-1))\\ &+\phi^{\rm P}_{\rm M_2}(x_2)(\phi^{\rm t}_{\rm V}(x_3)(x_1+x_2+x_3-1)\\ &-\phi^{\rm s}_{\rm V}(x_3)(x_1+x_2-x_3+3)))\big]\\ &\times h_{\rm anf1}(\alpha,\sqrt{|\beta_1^2|},b_1,b_2)E_{\rm anf}(t_{\rm c})\\ &-\big[\phi^{\rm A}_{\rm M_2}(x_2)\phi_{\rm V}(x_3)(x_3-1)(2r_{\rm V}^2-1)\\ &+r_{02}r_{\rm V}\big(\phi^{\rm P}_{\rm M_2}(x_2)(\phi^{\rm s}_{\rm V}(x_3)(x_1-x_2+x_3-1)\\ &+\phi^{\rm t}_{\rm V}(x_3)(x_1-x_2-x_3+1))\\ &+\phi^{\rm T}_{\rm M_2}(x_2)(\phi^{\rm s}_{\rm V}(x_3)(x_2+x_3-x_1-1)\\ &+\phi^{\rm t}_{\rm V}(x_3)(1-x_1+x_2-x_3)))\big]\\ &\times h_{\rm anf2}(\alpha,\sqrt{|\beta_2^2|},b_1,b_2)E_{\rm anf}(t_{\rm d})\Big\},\end{aligned}\tag{13}$$

with $r_{\rm V}=r_3=m_{\rm V}/m_{\rm D}$. For D→VP decays, the amplitudes are

$$\begin{aligned}\mathcal{A}^{\rm VP}_{\rm af} =& 8C_{\rm F}f_{\rm D}\pi m_{\rm D}^4\int_0^1{\rm d}x_2{\rm d}x_3\int_0^{1/\Lambda}b_2{\rm d}b_2b_3{\rm d}b_3\\ &\times\big\{\big[2\phi^{\rm s}_{\rm V}(x_2)r_{03}r_{\rm V}(\phi^{\rm T}_{\rm M_3}(x_3)x_3-\phi^{\rm P}_{\rm M_3}(x_3)(x_3-2))\\ &+\phi^{\rm A}_{\rm M_3}(x_3)\phi_{\rm V}(x_2)((2x_3-1)r_{\rm V}^2-x_3+1)\big]\\ &\times h_{\rm af}(\alpha,\beta,b_2,b_3)E_{\rm af}(t_{\rm a})\\ &-\big[2\phi^{\rm P}_{\rm M_3}(x_3)r_{03}r_{\rm V}(\phi^{\rm t}_{\rm V}(x_2)(x_2-1)\\ &+\phi^{\rm s}_{\rm V}(x_2)(x_2+1))-\phi^{\rm A}_{\rm M_3}(x_3)\phi_{\rm V}(x_2)(r_{\rm V}^2-1)x_2\big]\\ &\times h_{\rm af}(\alpha',\beta,b_3,b_2)E_{\rm af}(t_{\rm b})\big\},\end{aligned}\tag{14}$$

$$\begin{aligned}
\mathcal{M}_{\rm anf}^{\rm VP} = & \; 16\sqrt{\frac{2}{3}}C_{\rm F}\pi m_{\rm D}^4\int_0^1 {\rm d}x_1{\rm d}x_2{\rm d}x_3\int_0^{1/\Lambda} b_1{\rm d}b_1 b_2{\rm d}b_2 \\
& \times\phi_{\rm D}(x_1,b_1)\big\{\big[\phi_{\rm M_3}^{\rm A}(x_3)\phi_{\rm V}(x_2) \\
& \times(x_1+x_2+(-2x_1-2x_2+1)r_{\rm V}^2)) \\
& +r_{03}r_{\rm V}\left(\phi_{\rm M_3}^{\rm T}(x_3)(\phi_{\rm V}^{\rm s}(x_2)(1-x_1-x_2-x_3)\right. \\
& +\phi_{\rm V}^{\rm t}(x_2)(1-x_1-x_2+x_3)) \\
& +\phi_{\rm M_3}^{\rm P}(x_3)(\phi_{\rm V}^{\rm t}(x_2)(x_1+x_2+x_3-1) \\
& \left.+\phi_{\rm V}^{\rm s}(x_2)(x_1+x_2-x_3+3)))\right] \\
& \times h_{\rm anf1}(\alpha,\sqrt{|\beta_1^2|},b_1,b_2)E_{\rm anf}(t_{\rm c}) \\
& -\big[\phi_{\rm M_3}^{\rm A}(x_3)\phi_{\rm V}(x_2) \\
& \times(1-x_3+r_{\rm V}^2(x_1-x_2+2x_3-2)) \\
& +r_{03}r_{\rm V}\left(\phi_{\rm M_3}^{\rm P}(x_3)(\phi_{\rm V}^{\rm s}(x_2)(1-x_1+x_2-x_3)\right. \\
& +\phi_{\rm V}^{\rm t}(x_2)(x_1-x_2-x_3+1)) \\
& +\phi_{\rm M_3}^{\rm T}(x_3)(\phi_{\rm V}^{\rm s}(x_2)(x_2+x_3-x_1-1) \\
& \left.+\phi_{\rm V}^{\rm t}(x_2)(x_1-x_2+x_3-1)))\right] \\
& \times h_{\rm anf2}(\alpha,\sqrt{|\beta_2^2|},b_1,b_2)E_{\rm anf}(t_{\rm d})\big\}, \qquad (15)
\end{aligned}$$

with $r_{\rm V}=r_2=m_{\rm V}/m_{\rm D}$. The form of the wave functions of final state pseudoscalar mesons and vector mesons can be found in Ref. [13], with the different Gegenbauer moments used in this work as

$$\begin{aligned}
& a_{2\pi}^{\rm A}=0.70,\ a_{4\pi}^{\rm A}=0.45,\ a_{2\pi}^{\rm P}=0.70,\ a_{4\pi}^{\rm P}=0.36, \\
& a_{3\pi}^{\rm T}=0.80,\ a_{1\rm K}^{\rm A}=0.60,\ a_{2\rm K}^{\rm A}=0.10,\ a_{2\rm K}^{\rm P}=0.5, \\
& a_{4\rm K}^{\rm P}=-0.2,\ a_{3\rm K}^{\rm T}=0.65,\ a_{2\rho}^{\parallel}=a_{2\omega}^{\parallel}=0.6, \\
& a_{2\phi}^{\parallel}=0.70,\ a_{1\rm K^*}^{\parallel}=0.6,\ a_{2\rm K^*}^{\parallel}=0.11. \qquad (16)
\end{aligned}$$

Since the energy release in D decays is smaller than that in B decays, our light meson wave functions have larger $SU(3)$ breakings in D decays. For the distribution amplitudes of $\rm D/D_s$ meson, we take the same model as the B meson [13] with different hadronic parameter $\omega=0.35/0.5$ for $\rm D/D_s$ meson.

With the functions obtained in the above, the amplitudes of these pure annihilation decay channels can be given by

$$\begin{aligned}
& \mathcal{A}({\rm D^0\to K^{(*)0}\bar{K}^{(*)0}}) \\
= & \frac{G_{\rm F}}{\sqrt{2}}\Big\{V_{\rm cd}^*V_{\rm ud}\left[a_2\mathcal{A}_{\rm af}^{\rm K^{(*)0}\bar{K}^{(*)0}}+C_2\mathcal{M}_{\rm anf}^{\rm K^{(*)0}\bar{K}^{(*)0}}\right] \\
& +V_{\rm cs}^*V_{\rm us}\left[a_2\mathcal{A}_{\rm af}^{\rm \bar{K}^{(*)0}K^{(*)0}}+C_2\mathcal{M}_{\rm anf}^{\rm \bar{K}^{(*)0}K^{(*)0}}\right]\Big\}, \qquad (17)
\end{aligned}$$

$$\mathcal{A}({\rm D^0\to K^0\phi})=\frac{G_{\rm F}}{\sqrt{2}}V_{\rm cd}^*V_{\rm us}[a_2\mathcal{A}_{\rm af}^{\rm K\phi}+C_2\mathcal{M}_{\rm anf}^{\rm K\phi}], \qquad (18)$$

$$\mathcal{A}({\rm D^0\to \bar{K}^0\phi})=\frac{G_{\rm F}}{\sqrt{2}}V_{\rm cs}^*V_{\rm ud}[a_2\mathcal{A}_{\rm af}^{\rm \phi\bar{K}}+C_2\mathcal{M}_{\rm anf}^{\rm \phi\bar{K}}], \qquad (19)$$

$$\mathcal{A}({\rm D^+\to K^+\phi})=\frac{G_{\rm F}}{\sqrt{2}}V_{\rm cd}^*V_{\rm us}[a_1\mathcal{A}_{\rm af}^{\rm K\phi}+C_1\mathcal{M}_{\rm anf}^{\rm K\phi}], \qquad (20)$$

$$\begin{aligned}
\mathcal{A}({\rm D_s\to\pi^+\pi^0}) = & \frac{G_{\rm F}}{2}V_{\rm cs}^*V_{\rm ud}[a_2(\mathcal{A}_{\rm af}^{\pi^0\pi^+}-\mathcal{A}_{\rm af}^{\pi^+\pi^0}) \\
& +C_2(\mathcal{M}_{\rm anf}^{\pi^0\pi^+}-\mathcal{M}_{\rm anf}^{\pi^+\pi^0})] \\
& \sim 0, \qquad (21)
\end{aligned}$$

$$\begin{aligned}
\mathcal{A}({\rm D_s\to\pi^0\rho^+}) = & \frac{G_{\rm F}}{2}V_{\rm cs}^*V_{\rm ud}[a_2(\mathcal{A}_{\rm af}^{\pi^0\rho^+}-\mathcal{A}_{\rm af}^{\rho^+\pi^0}) \\
& +C_2(\mathcal{M}_{\rm anf}^{\pi^0\rho^+}-\mathcal{M}_{\rm anf}^{\rho^+\pi^0})], \qquad (22)
\end{aligned}$$

$$\begin{aligned}
\mathcal{A}({\rm D_s\to\pi^+\rho^0(\omega)}) = & \frac{G_{\rm F}}{\sqrt{2}}V_{\rm cs}^*V_{\rm ud}[a_2(\mathcal{A}_{\rm af}^{\pi^+\rho^0(\omega)} \\
& \mp\mathcal{A}_{\rm af}^{\rho^0(\omega)\pi^+})+C_2(\mathcal{M}_{\rm anf}^{\pi^+\rho^0(\omega)} \\
& \mp\mathcal{M}_{\rm anf}^{\rho^0(\omega)\pi^+})]. \qquad (23)
\end{aligned}$$

3 Numerical results and discussions

For numerical analysis, we use the following input parameters:

$$\begin{aligned}
& f_{\rm D/D_s}=0.23/0.257\ {\rm GeV},\ f_{\rm K}=0.16\ {\rm GeV},\ f_\pi=0.13\ {\rm GeV}, \\
& f_\rho^{\rm (T)}=0.209(0.165)\ {\rm GeV},\ f_{\rm K^*}^{\rm (T)}=0.217(0.185)\ {\rm GeV}, \\
& f_\omega^{\rm (T)}=0.195(0.145)\ {\rm GeV},\ f_\phi^{\rm (T)}=0.220(0.185)\ {\rm GeV}, \\
& |V_{\rm cd}|=0.2252\pm0.00065,\ |V_{\rm ud}|=0.9742\pm0.0002, \\
& |V_{\rm cs}|=0.97344\pm0.00016,\ |V_{\rm us}|=0.2253\pm0.00065, \\
& m_{0\pi}=1.4\ {\rm GeV},\ m_{0\rm K}=1.6\ {\rm GeV},\ \Lambda_{\rm QCD}^{\rm f=3}=0.375\ {\rm GeV}. \qquad (24)
\end{aligned}$$

After numerical calculation, the branching ratios of these decays together with experimental measurements [21] are listed in Table 1. We also list the results from the diagrammatic approach [4] and the pole model [7] for comparison.

The branching ratio obtained from the analytic formulas may be sensitive to many parameters especially those in the meson wave function. The theoretical uncertainties in our calculations, shown in Table 1, are caused by the variation of (i) the hadronic parameters, such as the shape parameters and the Gegenbauer moments in wave functions of initial and final state mesons; (ii) the unknown next-to-leading order QCD corrections and nonperturbative power corrections, characterized by the choice of the $\Lambda_{\rm QCD}=(0.375\pm0.05)$ GeV and the variations of the factorization scales defined in Eq. (5) and Eq. (9), respectively.

Table 1. Branching ratios (10^{-3}) for $D\to PP(V)$ decays together with experimental data [21], the recent results from the diagrammatic approach [4] and the predictions from the pole model [7].

decay modes	this work	Br(diagrammatic)	Br(pole model)	Br(Exp)
$D^0\to K^0\bar{K}^0$	$0.27^{+0.09}_{-0.08}$	0	0.3 ± 0.1	0.34 ± 0.08
$D_s\to\pi^+\pi^0$	0	0	0	<0.34
$D^0\to\bar{K}^0\phi$	$8.55^{+3.60}_{-3.41}$	8.68 ± 0.139	0.8 ± 0.2	8.34 ± 0.65
$D^0\to\bar{K}^0K^{*0}$	$0.44^{+0.20}_{-0.17}$	0.29 ± 0.22	0.16 ± 0.05	<0.56
$D^0\to K^0\bar{K}^{*0}$	$0.54^{+0.20}_{-0.15}$	0.29 ± 0.22	0.16 ± 0.05	<1.0
$D^0\to K^0\phi$	$0.012^{+0.004}_{-0.004}$	0.006 ± 0.005	0.020 ± 0.006	
$D^+\to K^+\phi$	$0.025^{+0.012}_{-0.008}$		0.020 ± 0.0020	
$D_s^+\to\pi^+\rho^0$	$2.11^{+0.87}_{-0.25}$		4.0 ± 4.0	0.2 ± 0.12
$D_s^+\to\pi^+\omega$	$0.050^{+0.029}_{-0.025}$		0	2.5 ± 0.7
$D_s^+\to\pi^0\rho^+$	$2.11^{+0.87}_{-0.24}$		4.0 ± 4.0	

In hadronic D decays, the $SU(3)$ breaking effect is remarkable, which can be demonstrated by the decay channel $D^0\to K^0\bar{K}^0$, with large branching ratio from experimental measurement. There are two kinds contributions from the quark pair $d\bar{d}$ and $s\bar{s}$ produced through a weak vertex. In $SU(3)$ limit, the two contributions exactly cancel each other due to the cancelation of the CKM matrix elements. Thus the diagrammatic approach [4] results in zero branching ratio for this channel. Taking the $SU(3)$ breaking effect into account, we give the result in agreement with the experimental data. For the decay $D^0\to\bar{K}^0\phi$, we reproduce the result of Ref. [9], which agrees well with the experimental data. For $D_s^+\to\pi^+\pi^0$ decay, the branching ratio vanishes due to the exact cancelation of the contributions from $u\bar{u}$ and $d\bar{d}$ components. In fact, this decay is forbidden because the two pions cannot form an s wave isospin 1 state due to the Bose-Einstein statistics. Any non-zero data for this decay may indicate the signal of new physics beyond the standard model.

For $D_s^+\to\pi^+\rho^0$ decay, the branching ratio is larger than the experimental result, while for $D_s^+\to\pi^+\omega$, it is much smaller than the experimental result. The reason is that the minus sign of ρ^0 relative to ω is compensated by the asymmetric space wave function of the two final states which are in the P-wave state. One possible solution is the soft final-state interactions as discussed in Ref. [4]. In general, the soft final-state interaction should be important in D meson decays, because there are many resonance states near the D meson mass, which may give severe pollution to D decays calculation. We expect more accurate measurements from experiments such as LHCb and BESS-III, which can help us understand better the dynamics of D meson decays.

4 Summary

In this work, we calculate the branching ratio of the 10 pure annihilation type $D_{(s)}\to PP(V)$ decays in the perturbative QCD factorization approach without considering soft final-states interactions. For most channels, our results agree well with the experimental data. The $SU(3)$ breaking effect is found to be remarkable, which can be indicated by the large branching ratio of $D^0\to K^0\bar{K}^0$ decay. We hope that the super B factories and BES-III can provide more accurate measurements for these decays, which will help us learn about the QCD dynamics in D meson decays and the annihilation mechanism.

We are very grateful to Yu Xin and Yu Fu-Sheng for helpful discussions.

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