Simultaneous Solution to $B_d^0 \to \phi K^{*0}$ Polarization Anomaly and $B_d \to \phi K_S$ *CP* Asymmetry*

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Abstract Based on the low-energy effective Hamiltonian with generalized QCD factorization, we calculate the new physics contributions to $B_d^0 \to \phi K^{*0}$ polarization anomaly and $B_d \to \phi K_S$ *CP* asymmetry induced by the neutral Higgs bosons H^0, h^0 and A^0 of the two Higgs doublet model Ⅲ. Within reasonable ranges of parameters, simultaneous solutions to the anomalies of the two processes give out $147 < |\lambda_{bs}\lambda_{ss}| < 165$. It will be really interesting in searching for the signs of new physics.

Key words two Higgs doublet model Ⅲ, QCD factorization, flavor changing neutral currents

1 Introduction

Looking for the signals of new physics (NP) beyond the Standard Model (SM) is one of the most important missions of high energy physics. The study of B physics is of great importance for testing indirect signals of new physics. In this respect, SLAC and KEK are doing a commendable job by providing us with a huge amount of data on various B-meson decays. Most of the experimental measurements are in perfect agreement with the SM. However, there still exist some puzzles observed in exclusive B-meson decays. Such as the abnormally small longitudinal polarizations fractions in $B_d^0 \to \phi K^{*0}$ decay[1—6] and the anomaly in the time-dependent *CP* asymmetry $S_{\phi K_S}$ for $B_d \to \phi K_S$ decay[7—10]. Confronted with these anomalies, we are forced not only to consider more precise QCD effects, but also to speculate on the existence of possible new physics.

The recent experimental results for the longitudinal polarizations fractions in $B_d^0 \to \phi K^{*0}$ decay are given as

$$\begin{aligned} f_L &= 0.52 \pm 0.05 \pm 0.02 \quad \text{BARBAR}^{[1]}, \\ f_L &= 0.45 \pm 0.05 \pm 0.02 \quad \text{Belle}^{[2]}, \\ f_L &= 0.57 \pm 0.10 \pm 0.05 \quad \text{CDF}^{[3]}. \end{aligned} \tag{1}$$

The *CP* asymmetry for the $B_d \to \phi K_S$ decay is given as

$$\begin{aligned} S_{\phi K_S} &= 0.19 \pm 0.32, \\ A_{\phi K_S} &= 0.12 \pm 0.20 \ \text{Belle}^{[7]}, \end{aligned} \tag{2}$$

$$\begin{aligned} S_{\phi K_S} &= 0.50 \pm 0.25^{+0.07}_{-0.04}, \\ A_{\phi K_S} &= 0.00 \pm 0.23 \pm 0.05 \ \text{BaBar}^{[8]}. \end{aligned} \tag{3}$$

It is well known that the flavor changing neutral current (FCNC) processes arise only from loop effects within the SM. Since $B_d^0 \to \phi K^{*0}$ and $B_d \to \phi K_S$ decays are the $\bar{b} \to \bar{s}s\bar{s}$ pure penguin process, they could be used as effective probes of NP scenarios. Especially, using PQCD approach, the abnormally small longitudinal polarizations in $B_d^0 \to \phi K^{*0}$ decay in the two Higgs doublet model (2HDM) Ⅲ have been studied[4]. Moreover, the solution to the

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anomaly in the time-dependent CP asymmetry $S_{\phi K_S}$ for $B_d \to \phi K_S$ decay is also discussed by Ref. [10]. In this paper, we use QCD factorization approach[6] and try to provide simultaneous solutions to the abnormally small longitudinal polarizations in $B_d^0 \to \phi K^{*0}$ decay and the anomaly in the time-dependent CP asymmetry $S_{\phi K_S}$ for $B_d \to \phi K_S$ decay within the reasonable parameter space of the 2HDM Ⅲ.

2 The general characteristics of 2HDM Ⅲ

One of the most popular extensions of the SM is the so-called two Higgs doublet model, which has two complex Higgs doublets instead of only one in the SM. The cases of the same or the two different Higgs fields to up- and down-type quarks are called model Ⅰ or model Ⅱ. While in 2HDM Ⅲ, both the doublets can couple to the up- and down-type quarks[11, 12].

Generally we can write down Yukawa Lagrangian for model Ⅲ[11—13]

$$\mathscr{L}_{\rm Y} = \eta_{ij}^U \bar{Q}_{i,\rm L}\tilde{\phi}_1 U_{j,\rm R} + \eta_{ij}^D \bar{Q}_{i,\rm L}\phi_1 D_{j,\rm R} + \xi_{ij}^U \bar{Q}_{i,\rm L}\tilde{\phi}_2 U_{j,\rm R} + \xi_{ij}^D \bar{Q}_{i,\rm L}\phi_2 D_{j,\rm R} + {\rm h.c.}\ , \tag{4}$$

where $\phi_i (i=1,2,)$ denote the two scalar doublets, $\tilde{\phi}_{1,2} = {\rm i}\tau_2\phi_{1,2}^*$, $Q_{i,\rm L}$ with $i=(1,2,3)$ is the left-handed fermion doublet, $U_{j,\rm R}$ and $D_{j,\rm R}$ are the right-handed singlets. $Q_{i,\rm L}, U_{j,\rm R}$ and $D_{j,\rm R}$ are all weak eigenstates, which can be rotated into mass eigenstates. While $\eta^{U,D}$ and $\xi^{U,D}$ are generally the non-diagonal matrices of the Yukawa coupling.

One can express ϕ_1 and ϕ_2 in a suitable basis such that only the first doublet generates all the gauge-boson and fermion masses, i.e., such that

$$\langle\phi_1\rangle = \begin{pmatrix} 0 \\ \frac{v}{\sqrt{2}} \end{pmatrix}, \quad \langle\phi_2\rangle = 0\ . \tag{5}$$

The two doublets in the basis then have the form

$$\phi_1 = \frac{1}{\sqrt{2}}\left[\begin{pmatrix} 0 \\ v+\chi_1^0 \end{pmatrix} + \begin{pmatrix} \sqrt{2}G^+ \\ {\rm i}G^0 \end{pmatrix}\right], \quad \phi_2 = \frac{1}{\sqrt{2}}\begin{pmatrix} \sqrt{2}H^+ \\ \chi_2^0 + {\rm i}A^0 \end{pmatrix}, \tag{6}$$

where $G^{0,\pm}$ are the Goldstone bosons that would be eaten away in the Higgs mechanism to become the longitudinal components of the weak gauge bosons. $H^\pm$ and A^0 are the physical charged Higgs boson and CP-odd neutral Higgs boson, respectively. χ_1^0 and χ_2^0 are not the neutral mass eigenstates but linear combinations of the CP-even neutral Higgs boson mass eigenstates H^0 and h^0

$$\chi_1^0 = H^0\cos\alpha - h^0\sin\alpha, \tag{7}$$

$$\chi_2^0 = H^0\sin\alpha + h^0\cos\alpha, \tag{8}$$

where α is the mixing angle, the advantage of choosing the basis is no couplings of the form $\chi_2^0 ZZ$ and $\chi_2^0 W^+W^-$.

After the diagonalization of mass matrix of the fermions fields, the Yukawa Lagrangian becomes

$$\mathscr{L}_{\rm YC} = \hat{\eta}_{ij}^U \bar{Q}_{i,\rm L}\tilde{\phi}_1 U_{j,\rm R} + \hat{\eta}_{ij}^D \bar{Q}_{i,\rm L}\phi_1 D_{j,\rm R} + \hat{\xi}_{ij}^U \bar{Q}_{i,\rm L}\tilde{\phi}_2 U_{j,\rm R} + \hat{\xi}_{ij}^D \bar{Q}_{i,\rm L}\phi_2 D_{j,\rm R} + {\rm h.c.} \tag{9}$$

where $Q_{i,\rm L}$, $U_{j,\rm R}$, and $D_{j,\rm R}$ denote the fermion fields which are in mass eigenstates and

$$\hat{\eta}^{U,D} = (V_{\rm L}^{U,D})^{-1}\cdot\eta^{U,D}\cdot V_{\rm R}^{U,D} = \frac{\sqrt{2}}{v}M^{U,D}, \tag{10}$$

$$\hat{\xi}^{U,D} = (V_{\rm L}^{U,D})^{-1}\cdot\xi^{U,D}\cdot V_{\rm R}^{U,D}\ . \tag{11}$$

$V_{\rm L,R}^{U,D}$ are the rotation matrices acting on the up- and down-type quarks, with left or right chirality, respectively. $V_{\rm CKM} = (V_{\rm L}^U)^\dagger V_{\rm L}^D$ is the usual Cabibbo-Kobayashi-Maskawa (CKM) matrix. From Eq. (9), one can see that the matrices $\hat{\xi}^{U,D}$ allow scalar-mediated FCNC, and in the quark mass basis only the matrices $\hat{\eta}^{U,D}$ are diagonal, but the matrices $\hat{\xi}^{U,D}$ are in general not diagonal.

The matrix $\hat{\xi}^{U,D}$ contain the FCNC couplings. In this paper, we use the Cheng-Sher relation[13]

$$\xi_{ij}^{U,D} = \lambda_{ij}\frac{\sqrt{m_i m_j}}{v}\ , \tag{12}$$

by which the quark-mass hierarchy ensures that the FCNC within the first two generations are naturally suppressed by the small quark masses, while a large freedom is allowed for the FCNC involving the third generations.

3 Polarization anomaly of $B_d^0 \to \phi K^{*0}$ decay

In SM, the analysis of the $B_d^0 \to \phi K^{*0}$ decay can be performed in terms of an effective low energy theory with the Hamiltonian[14]

$$H_{\rm eff}^{\rm SM} = -\frac{G_{\rm F}}{\sqrt{2}} V_{\rm tb} V_{\rm ts}^* \sum_{i=3}^{10} C_i O_i \ . \tag{13}$$

The amplitude for the decay can be written as[6]

$$A_\lambda = -\frac{G_{\rm F}}{\sqrt{2}} V_{\rm tb} V_{\rm ts}^* a_\lambda(\phi K^{*0}) X^{(BK^{*0},\phi)} \ , \tag{14}$$

where

$$\begin{aligned} X^{(BK^{*0},\phi)} = & \langle\phi|\bar{s}\gamma_\mu s|0\rangle\langle K^{*0}|\bar{s}\gamma^\mu(1-\gamma_5)b|B\rangle = \\ & {\rm i} f_\phi M_\phi \left[\varepsilon_1^* \cdot \varepsilon_2^* (M_{\rm B}+M_{\rm K^{*0}}) A_1(M_\phi^2) - \right. \\ & (\varepsilon_1^* \cdot P_{\rm B})(\varepsilon_2^* \cdot P_{\rm B}) \frac{2A_2(M_\phi^2)}{M_{\rm B}+M_{\rm K^{*0}}} + \\ & \left. {\rm i}\epsilon_{\mu\nu\alpha\beta}\varepsilon_2^{*\mu}\varepsilon_1^{*\nu} P_{\rm B}^\alpha P_{\rm K^{*0}}^\sigma \frac{2V(M_\phi^2)}{M_{\rm B}+M_{\rm K^{*0}}}\right], \end{aligned} \tag{15}$$

and $a_\lambda(\phi K^{*0})$ is the effective coefficient which depends on the factorization approach.

Since B meson is a pseudoscalar, the final two vector mesons must have the same helicity. In the helicity basis, taking the $\phi(K^{*0})$ meson flying in the minus (plus) z-direction and using the sign convention $\epsilon^{0123} = -1$, the amplitude can be decomposed into three helicity amplitudes

$$\begin{aligned} H_0 = & \frac{G_{\rm F}}{\sqrt{2}} V_{\rm tb} V_{\rm ts}^* a_0(\phi K^{*0}) \frac{{\rm i} f_\phi}{2M_{\rm K^{*0}}} \times \\ & \left[(M_{\rm B}^2 - M_{\rm K^{*0}}^2 - M_\phi^2)(M_{\rm B}+M_{\rm K^{*0}}) A_1(M_\phi^2) - \right. \\ & \left. \frac{4M_{\rm B}^2 p_{\rm c}^2 A_2(M_\phi^2)}{M_{\rm B}+M_{\rm K^{*0}}}\right], \end{aligned} \tag{16}$$

$$\begin{aligned} H_\pm = & {\rm i}\frac{G_{\rm F}}{\sqrt{2}} V_{\rm tb} V_{\rm ts}^* a_\pm(\phi K^{*0}) M_\phi f_\phi \times \\ & \left[(M_{\rm B}+M_{\rm K^{*0}}) A_1(M_\phi^2) \mp \frac{2M_{\rm B} p_{\rm c}}{M_{\rm B}+M_{\rm K^{*0}}} V(M_\phi^2)\right], \end{aligned}$$

where

$$p_{\rm c} = \frac{\sqrt{[M_{\rm B}^2-(M_{\rm K^{*0}}+M_\phi)^2][M_{\rm B}^2-(M_{\rm K^{*0}}-M_\phi)^2]}}{2M_{\rm B}}, \tag{17}$$

denotes the center of mass momentum of the meson ϕ or K^{*0} in the B_d^0 rest frame.

Then the branching ratio is written as

$$B_{\rm r} = \frac{\tau_{\rm B} p_{\rm c}}{8\pi M_{\rm B}^2}(|H_0|^2+|H_+|^2+|H_-|^2) \ , \tag{18}$$

and the longitudinal and the transverse polarization rates are

$$\begin{aligned} f_{\rm L} = & \frac{\Gamma_{\rm L}}{\Gamma} = \frac{|H_0|^2}{|H_0|^2+|H_+|^2+|H_-|^2}, \\ f_{\rm T} = & \frac{\Gamma_{\rm T}}{\Gamma} = \frac{|H_+|^2+|H_-|^2}{|H_0|^2+|H_+|^2+|H_-|^2} \ . \end{aligned} \tag{19}$$

FCNC processes could happen at one loop level in the SM, but are suppressed by the GIM mechanism. The discrepancy between the SM result and the B factory data indicates the presence of NP which may enter through new FCNC processes arising at the tree level.

In this paper, we only consider fermions couplings with the neutral gauge bosons H^0, h^0, and A^0 in the 2HDM Ⅲ. The FCNC part of the Yukawa Lagrangian could be[11, 15]

$$\begin{aligned} \mathscr{L}_{\rm Y,FCNC} = & -\frac{H^0\sin\alpha + h^0\cos\alpha}{\sqrt{2}} \times \\ & \left\{\bar{D}\left[\hat{\xi}^D \frac{1}{2}(1+\gamma_5) + \hat{\xi}^{D\dagger}\frac{1}{2}(1-\gamma_5)\right] D\right\} - \\ & \frac{{\rm i}A^0}{\sqrt{2}}\left\{\bar{D}\left[\hat{\xi}^D \frac{1}{2}(1+\gamma_5) - \hat{\xi}^{D\dagger}\frac{1}{2}(1-\gamma_5)\right] D\right\} - \\ & \frac{g}{2M_{\rm W}}(H^0\cos\alpha - h^0\sin\alpha)\bar{D} M_{\rm D} D, \end{aligned} \tag{20}$$

where α is the mixing angle, D represents the mass eigenstates of down-type quarks. We only consider s quark in this paper. ξ^D is defined by Eq. (11).

The effective Hamiltonians of the $\bar{b} \to \bar{s}s\bar{s}$ transition mediated by the H^0, h^0, and A^0 are given by, respectively

$$\begin{aligned} H_{\rm eff}^{H^0} = & -\frac{\sqrt{2}G_{\rm F}}{2M_{H^0}^2}\sin^2\alpha\, m_{\rm s}\sqrt{m_{\rm b} m_{\rm s}} \times \\ & \left[\lambda_{\rm ss}\lambda_{\rm bs}^*(O_7+O_9) + \lambda_{\rm ss}^*\lambda_{\rm bs}(O_7'+O_9')\right], \end{aligned} \tag{21}$$

$$\begin{aligned} H_{\rm eff}^{h^0} = & -\frac{\sqrt{2}G_{\rm F}}{2M_{h^0}^2}\cos^2\alpha\, m_{\rm s}\sqrt{m_{\rm b} m_{\rm s}} \times \\ & \left[\lambda_{\rm ss}\lambda_{\rm bs}^*(O_7+O_9) + \lambda_{\rm ss}^*\lambda_{\rm bs}(O_7'+O_9')\right], \end{aligned} \tag{22}$$

$$\begin{aligned} H_{\rm eff}^{A^0} = & -\frac{\sqrt{2}G_{\rm F}}{2M_{A^0}^2} m_{\rm s}\sqrt{m_{\rm b} m_{\rm s}} \times \\ & \left[\lambda_{\rm ss}\lambda_{\rm bs}^*(O_7+O_9) + \lambda_{\rm ss}^*\lambda_{\rm bs}(O_7'+O_9')\right], \end{aligned} \tag{23}$$

where the constant term, and the coefficients of $\lambda_{ss}(\lambda_{ss}^*)$, $\lambda_{bs}(\lambda_{bs}^*)$ terms which are 100 times smaller than the coefficients of $\lambda_{bs}\lambda_{bs}^*$, $\lambda_{bs}^*\lambda_{bs}$, so we have neglected them in the upper Eqs. (20—22). $H_{eff}^{H^0}$, $H_{eff}^{h^0}$, and $H_{eff}^{A^0}$ have the same operators O_7 and O_9 as in the SM effective Hamiltonian, O_7' and O_9' are new operators

$$O_7' = (\bar{q}q)_{V-A}(\bar{b}s)_{V+A},\quad O_9' = (\bar{q}q)_{V+A}(\bar{b}s)_{V+A}, \quad (24)$$

the hadronic matrix elements of these operators are calculated up to the α_s order using the QCD factorization approach in this paper. For convenience, we assume that λ_{ss} is real, and $\lambda_{bs} = |\lambda_{bs}|e^{i\phi}$.

The additional contributions to the SM Wilson coefficients at the M_W scale are

$$\begin{aligned}\Delta C_7 \;=\; & \Delta C_9 = m_s\sqrt{m_b m_s}\frac{|V_{tb}V_{ts}^*|}{V_{tb}V_{ts}^*}\times \\ & \left(\frac{\sin^2\alpha}{M_{H^0}^2}+\frac{\cos^2\alpha}{M_{h^0}^2}+\frac{1}{M_{A^0}^2}\right)|\lambda_{bs}^*\lambda_{ss}|e^{i\phi} = \\ & -m_s\sqrt{m_b m_s}\left(\frac{\sin^2\alpha}{M_{H^0}^2}+\frac{\cos^2\alpha}{M_{h^0}^2}+\frac{1}{M_{A^0}^2}\right)\times \\ & |\lambda_{bs}\lambda_{ss}|e^{-i\phi}\ . \end{aligned} \quad (25)$$

The new operators contribution to the Wilson coefficients at the M_W scale are

$$\begin{aligned}\Delta C_7' \;=\; & \Delta C_9' = m_s\sqrt{m_b m_s}\frac{|V_{tb}V_{ts}^*|}{V_{tb}V_{ts}^*}\times \\ & \left(\frac{\sin^2\alpha}{M_{H^0}^2}+\frac{\cos^2\alpha}{M_{h^0}^2}+\frac{1}{M_{A^0}^2}\right)|\lambda_{bs}\lambda_{ss}^*|e^{i\phi} = \\ & -m_s\sqrt{m_b m_s}\left(\frac{\sin^2\alpha}{M_{H^0}^2}+\frac{\cos^2\alpha}{M_{h^0}^2}+\frac{1}{M_{A^0}^2}\right)\times \\ & |\lambda_{bs}\lambda_{ss}|e^{i\phi}\ . \end{aligned} \quad (26)$$

We neglect the renormalization group running between M_W and M_{H^0}, M_{h^0}, M_{A^0}. The full description of the running of the Wilson coefficient from the M_W scale to m_b can be found in Ref. [14]. We only repeat the directly relevant steps. The renormalization group equation for the Wilson coefficients $\boldsymbol{C}$

$$\frac{d}{d\ \ln\mu}\boldsymbol{C} = \gamma^T(g)\boldsymbol{C}(\mu), \quad (27)$$

can be solved with the help of the U matrix which describes the QCD evolution

$$\boldsymbol{C}(\mu) = U(\mu, M_W)\boldsymbol{C}(M_W), \quad (28)$$

where $\gamma^T(g)$ is the transpose of the anomalous dimension matrix $\gamma(g)$. With the help of $dg/d\ln\mu = \beta(g)$, U obeys the same equation as $\boldsymbol{C}(\mu)$. As for the ΔC_9 and $\Delta C_9'$, one can get their anomalous dimensions to the first order in α_s

$$\gamma(\alpha_s) = \frac{\alpha_s}{4\pi}\gamma^{(0)} = \frac{\alpha_s}{4\pi}\begin{pmatrix}-6/N & 6\\ 6 & -6/N\end{pmatrix}, \quad (29)$$

and for the ΔC_7 and $\Delta C_7'$, we have the opposite anomalous dimension matrix.

We get the total decay amplitude

$$H_\lambda = H_\lambda^{SM} + H_\lambda^{new}, \quad (30)$$

where H_λ^{new} denotes the contributions from H^0, h^0, and A^0.

Taking into account the constrains on parameters from experimental data and theoretical limits, we take the NLO Wilson coefficients evaluated at the scale of $\mu = m_b$[14], the decay constants $f_B = 0.216$GeV, $f_\phi = 0.231$GeV and the form factors of light-cone QCD sum rules[16, 17]. The integral of B meson light-cone distribution amplitude can be parameterized as[18]

$$\int_0^1 d\xi\frac{\phi_B(\xi)}{\xi} = \frac{M_B}{\Lambda_B}, \quad (31)$$

and we take $\Lambda_B = 0.35$GeV in our calculations. The logarithmic and linear divergences may be phenomenologically parameterized as[18, 19]

$$\int_0^1\frac{dy}{y} = \ln\frac{M_B}{\Lambda_h},\quad \int_0^1\frac{dy}{y^2} = \frac{M_B}{\Lambda_h}, \quad (32)$$

with $\Lambda_h = 0.5$GeV.

The parameters in the 2HDM Ⅲ can be taken as $\alpha = \frac{\pi}{4}$, $M_{H^0} = 150$GeV, $M_{h^0} = 115$GeV, $M_{A^0} = 120$GeV[20]. Considering the NP contributions and combining the theoretical formulas and input parameters, the amplitude of $B_d^0 \to \phi K^{*0}$ decay depends only on two parameters $|\lambda_{bs}\lambda_{ss}|$ and ϕ. We can put constraint on the possible regions for the two NP parameters from the measured branching ratio and the longitudinal polarization of the $B_d^0 \to \phi K^{*0}$ decay. Considering the branching ratio at 2σ and the longitudinal polarization at 1σ, we give the contour plots for branching ratio and the longitudinal polarization having $|\lambda_{bs}\lambda_{ss}|$ as a function of the new phase ϕ in Fig. 1.

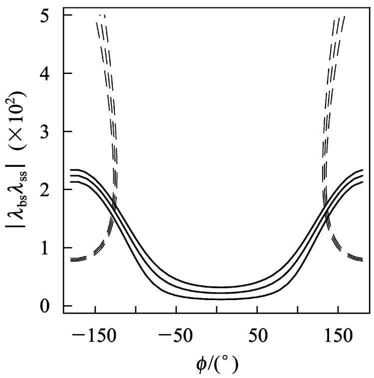

Fig. 1. The solid curves are the contour plots for $Br = (9.5 \pm 0.9) \times 10^{-6}$ within 2σ and the dashed curves for $f_L = 0.49 \pm 0.04$ within 1σ. The central curves are for the central values of Br and f_L, respectively.

From Fig. 1, the experimental values of the branching ratio and the longitudinal polarization for $B_d^0 \to \phi K^{*0}$ decay can be obtained simultaneously in the overlapped regions, which are

$$147 < |\lambda_{bs}\lambda_{ss}| < 176. \tag{33}$$

In Ref. [11], B_s^0–$\bar{B}_s^0$ mixing points out that the λ_{bs} coupling may not be small, it can be somewhat bigger than one, which is a lower bound. The λ_{ss} coupling can be as large as $O(100)$, we find that this parameter space is reasonable. If the recent experimental measurements on the polarization anomaly are credible, comparing the former constrains to the coupling parameters, we also obtain more stringent bound.

4 CP asymmetry for $B_d \to \phi K_S$ decay

In order to calculate the contributions on CP asymmetry for the $B_d \to \phi K_S$ decay, we take

$$S_{\phi K_S} = 0.40 \pm 0.20, \quad A_{\phi K_S} = 0.07 \pm 0.15, \tag{34}$$

which are the average value of Belle and BARBAR measurements in Eq. (2) and Eq. (3).

We adopt the QCD factorization for $B \to PV$[18], the decay amplitude $B_d \to \phi K_S$ is

$$\begin{aligned} A(B_d \to \phi K_S) = & \frac{G_F}{\sqrt{2}} \sum_{p=u,c} \lambda_p a(\phi K_S) \langle \phi | \bar{s}\gamma_\mu s | 0 \rangle \times \\ & \langle K_S | \bar{s}\gamma^\mu (1 \mp \gamma_5) b | B \rangle = \\ & -i\sqrt{2} G_F \sum_{p=u,c} \lambda_p a(\phi K_S) \times \\ & f_\phi M_\phi F_1^{BK_S}(M_\phi^2)(\varepsilon^* \cdot P_B), \end{aligned} \tag{35}$$

where $a(\phi K_S)$ is the effective coefficient.

The branching ratio of $B_d \to \phi K_S$ decay in the B meson rest frame can be written as

$$B_r = \frac{\tau_B p_c}{8\pi M_B^2} \left| A(B_d \longrightarrow \phi K_S) \right|^2, \tag{36}$$

where

$$p_c = \frac{\sqrt{[M_B^2 - (M_{K_S} + M_\phi)^2][M_B^2 - (M_{K_S} - M_\phi)^2]}}{2M_B}, \tag{37}$$

denotes the center of mass momentum of the meson ϕ or K_S in the B_d rest frame.

For the $B_d \to \phi K_S$ decay, the time-dependent CP asymmetry is

$$\begin{aligned} a_{\phi K_S} = & \frac{\Gamma(\bar{B}^0(t) \to \phi K_S) - \Gamma(B^0(t) \to \phi K_S)}{\Gamma(\bar{B}^0(t) \to \phi K_S) + \Gamma(B^0(t) \to \phi K_S)} = \\ & A_{\phi K_S} \cos(\Delta M_{B_d} t) + S_{\phi K_S} \sin(\Delta M_{B_d} t), \end{aligned} \tag{38}$$

where the direct and the indirect CP asymmetry parameters are given by, respectively

$$A_{\phi K_S} = \frac{|\lambda_{\phi K_S}|^2 - 1}{|\lambda_{\phi K_S}|^2 + 1}, \quad S_{\phi K_S} = \frac{2\mathrm{Im}[\lambda_{\phi K_S}]}{|\lambda_{\phi K_S}|^2 + 1}. \tag{39}$$

The parameter $\lambda_{\phi K_S}$ is defined by

$$\lambda_{\phi K_S} \equiv \eta_{\phi K_S} \frac{V_{tb}^* V_{td}}{V_{tb} V_{td}^*} \cdot \frac{V_{cs} V_{cd}^*}{V_{cs}^* V_{cd}} \cdot \frac{\bar{A}(\phi \bar{K}_S)}{A(\phi K_S)}, \tag{40}$$

where $\eta_{\phi K_S} = -1$ is CP eigenvalue of the ϕK_S state.

Considering the left-handed and right-handed flavor changing couplings from the contributions of H^0, h^0, and A^0, the additional contributions $\Delta C_{9(7)}$ and $\Delta C'_{9(7)}$ are the same as in the $B_d^0 \to \phi K^{*0}$ decay.

In the QCD factorization approach, we use the same input parameters as the process of $B_d^0 \to \phi K^{*0}$ decay, and get the $|\lambda_{bs}\lambda_{ss}| - \phi$ relations from the contour plots for the branching ratio and the CP asymmetry in Fig. 2.

Combining the constrains from the branching ratio for $B_d \to \phi K_S$ decay at 2σ and both $A_{\phi K_S}$ and $S_{\phi K_S}$ at 2σ, we find that the allowed regions are

$$\begin{aligned} 140 < |\lambda_{bs}\lambda_{ss}| < 165, \\ 120 < |\lambda_{bs}\lambda_{ss}| < 145, \\ 60 < |\lambda_{bs}\lambda_{ss}| < 85. \end{aligned} \tag{41}$$

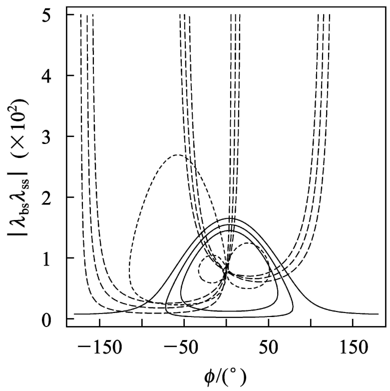

Fig. 2. The solid curves are the contour plots for the current world average value $Br = (8.3^{+1.2}_{-1.0})\times10^{-6}$ within 2σ[21], the dashed curves and dotted curves are for the current experimental results $S_{\phi K_S}=0.40\pm0.20$ and $A_{\phi K_S}=0.07\pm0.15$ within 2σ, respectively. The central curves are for the central values of branch ratio, $S_{\phi K_S}$ and $A_{\phi K_S}$, respectively.

5 Summary

In summary, in this paper we study the polarization anomaly of the $B_d^0\to\phi K^{*0}$ decay and the anomaly CP asymmetry $S_{\phi K_S}$ for $B_d\to\phi K_S$ decay in the 2HDM III. We consider the additional contributions from the neutral gauges boson H^0, h^0, and A^0, which induce the FCNCs at the tree level. The effective Hamiltonian $H_{eff}^{H^0}$, $H_{eff}^{h^0}$, and $H_{eff}^{A^0}$ include some new operators which do not exist in SM. As some operators in SM, the hadronic matrix elements of the new operators have been calculated up to the α_s order using the QCD factorization approach. We select $|\lambda_{bs}\lambda_{ss}|$ and the phase ϕ as the parameters and express the relation of the branching ratio and the longitudinal polarization, we give out the contour plots for branching ratio and the longitudinal polarization having $|\lambda_{bs}\lambda_{ss}|$ as a function of the new phase ϕ, and find that the large longitudinal polarization can be reduced to the experimental results within the reasonable parameters.

In the same way, we discuss the $B_d\to\phi K_S$ decay and show the images of $|\lambda_{bs}\lambda_{ss}|-\phi$ from the contour plots for the branching ratio and the CP asymmetry.

From Eqs. (33) and (41), we can get the self-consistent results of

$$147<|\lambda_{bs}\lambda_{ss}|<165 \qquad (42)$$

for the two processes in $B_d^0\to\phi K^{*0}$ and $B_d\to\phi K_S$. And the results we get naturally are very consistent with the constrains and the assumptions on the model parameters.

In conclusion, we have showed that the polarization anomaly in $B_d^0\to\phi K^{*0}$ decay and the anomaly in the time-dependent CP asymmetry $S_{\phi K_S}$ for $B_d\to\phi K_S$ decay, which are difficult to be explained in SM, but can be explained simultaneously in the 2HDM III. Moreover, if the recent experimental measurements on the polarization anomaly and CP asymmetry are credible, these limits on the coupling parameters that we give out will really be interesting in searching for the signs of new physics.

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$B_d^0 \to \phi K^{*0}$ 极化反常和 $B_d \to \phi K_S$ CP 不对称的研究*

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摘要 在模型Ⅲ下, 仅考虑来自于中性规范玻色子 H^0, h^0 和 A^0 产生的树图阶味改变中性流的贡献, 并利用QCD因子化方法, 对 $B_d^0 \to \phi K^{*0}$ 衰变过程的极化反常和 $B_d \to \phi K_S$ 衰变过程的 CP 不对称进行了研究, 经过计算发现, 在 $147 < |\lambda_{bs}\lambda_{ss}| < 165$ 的参数范围内, 这两个过程中存在的反常现象能够同时给予解释.

关键词 模型Ⅲ QCD因子化 味改变中性流

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