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Analysis of $B_d^0 \to \phi K^{*0}$ decay mode with supersymmetry

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Abstract. Motivated by the recent measurement of a low longitudinal polarization fraction in the decay mode $B_d^0 \to \phi K^{*0}$, which appears not to be in agreement with the standard model expectation, we analyze this mode in the minimal supersymmetric standard model with the mass insertion approximation. Within the standard model, with the factorization approximation, the longitudinal polarization is expected to be $f_{\rm L} \sim 1 - \mathcal{O}(1/m_b^2)$. We find that this anomaly can be explained in the minimal supersymmetric standard model with either the LR or the RL mass insertion approximation.

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1 Introduction

One of the important goals of the B-factories is to verify the standard model (SM) predictions and to serve as a potential avenue to reveal new physics beyond the SM. A huge collection of data in the B-sector has already been accumulated at both the B-factories (Belle and BABAR). This in turn has led, for the first time, to the observation of CP violation in the B system, outside the kaon system. In fact, the angle β of the unitarity triangle has been measured from the time dependent CP asymmetry of the gold plated $B_d^0 \to J/\psi K_{\rm S}$ mode by both Belle and BABAR, with almost similar values. The current world average of $\sin 2\beta$ is [1]

$$(\sin 2\beta)^{b\to c\bar{c}s} = 0.685 \pm 0.032\,, \tag{1}$$

which is consistent with the SM expectation. With the accumulation of more and more data the experimental B-physics is now all set to enter an unmatched precision era. Unfortunately, we have not seen any clear evidence of physics beyond the SM so far, as far as B-physics is concerned.

Already there are some more measurements available at the B-factories, which are not as clean as the $\sin 2\beta$ measurement in the golden decay mode $B \to J/\psi K_{\rm S}$, but from the pattern of deviation observed it appears that these measurements, in the long run with accumulation of more data, may reveal the signature of new physics. One of the modes of this kind is the decay mode $B \to \phi K_{\rm S}$, where in the SM one expects to obtain the same value of $\sin 2\beta$ from its CP asymmetry measurements, as in the case of $B \to J/\psi K_{\rm S}$, with a correction of $\mathcal{O}(\lambda^2)$ [2]. The basic difference between these two modes is that the golden mode is tree dominated ($b \to c\bar{c}s$) whereas the decay mode $B \to \phi K_{\rm S}$ is penguin dominated ($b \to s\bar{s}s$). It should be recalled here that in earlier times the deviation between these two measurements was very large but with the accumulation of more data the difference has reduced somewhat. The present averaged value is [1]

$$(\sin 2\beta)_{\phi K_{\rm S}} = 0.47 \pm 0.19\,, \tag{2}$$

which has about 1σ deviation from the corresponding $c\bar{c}$ measurements. In the future, even if the $(\sin 2\beta)_{\phi K_{\rm S}}$ value stabilizes around the present central value, with error bars reduced it might show the presence of new physics (NP). Moreover, there are other decay modes (involving the $b \to s\bar{s}s$ transition) where the data show a similar trend. Except for the decay mode $B^0 \to \eta' K^0$, the value of $\sin 2\beta$ extracted from all such modes are within the 1σ deviation from the corresponding $c\bar{c}$ value [3]. The present average value of the $B^0 \to \eta' K^0$ mode is $(\sin 2\beta)_{\eta' K^0} = 0.48 \pm 0.09$, which shows about 2.3σ deviation from (1).

The vector–vector counterpart of the seemingly problematic $B \to \phi K_{\rm S}$ decay mode, i.e., $B \to \phi K^*$, governed by the same $b \to s\bar{s}s$ transition as in $B \to \phi K_{\rm S}$, has also created a lot of attention recently. Both BABAR [4] and Belle [5] have observed this decay mode and the measured quantities are summarized in Table 1. The measured longitudinal polarization fraction in this mode is well below its expected value [6], i.e., $f_{\rm L} \sim 1 - \mathcal{O}(1/m_b^2)$, widely known in the literature as the polarization anomaly in $B \to \phi K^*$. In the B rest frame both the vector mesons are emitted back-to-back and from the spin angular momentum conservation it follows that both the vector mesons are paired up with the same helicity combinations (like 00, ++ and −−, i.e., the helicity combinations out of the three pos-

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Table 1. Experimental values of branching ratio (in units of 10^{-6}), polarization fractions and triple product asymmetries for the $B_d^0 \to \phi K^{*0}$ mode. The Belle results for the triple product asymmetries are obtained from combined ϕK^{*0} and ϕK^{*+} data

Observables	BABAR [4]	Belle [5]	Average
$\mathcal{B}$	$9.2 \pm 0.9 \pm 0.5$	$10.0^{+1.6+0.7}_{-1.5-0.8}$	9.4 ± 0.9
$f_{\rm L}$	$0.52 \pm 0.05 \pm 0.02$	$0.45 \pm 0.05 \pm 0.02$	0.49 ± 0.04
$f_\perp$	$0.22 \pm 0.05 \pm 0.02$	$0.30 \pm 0.06 \pm 0.02$	0.25 ± 0.04
$\mathcal{A}_{\rm T}^{\parallel}$	$-0.02 \pm 0.04 \pm 0.01$	$\frac{1}{2}[0.01 \pm 0.10 \pm 0.02]$	-0.01 ± 0.03
$\mathcal{A}_{\rm T}^{0}$	$0.11 \pm 0.05 \pm 0.01$	$\frac{1}{2}[0.16^{+0.16}_{-0.14} \pm 0.03]$	0.10 ± 0.04

sible helicity states for each vector meson, namely, $\lambda = 0, +$ and $-$). In the SM, it so happens that the helicity combination 00, called longitudinal, (i.e., for both the vector mesons the spin direction is proportional to the direction of motion) is almost the only preferred one and the occurrence of two other possible helicity combinations is suppressed by $\mathcal{O}(1/m_B^2)$, m_B being the B-meson mass. Thus, the longitudinal polarization fraction is defined as the ratio of the decay rate corresponding to the longitudinal polarization (say $\Gamma_{\rm L}$) to that of the total decay rate Γ, i.e., $f_{\rm L} = \Gamma_{\rm L}/\Gamma \approx 1$. However, as seen from Table 1 its measured value is only about $\approx 50\%$ of the expected value.

The unexpected deviation of $f_{\rm L}$ from the expected value of $\mathcal{O}(1)$ is known as the polarization anomaly in $B \to \phi K^*$ decay. In practice, the value of $f_{\rm L}$ is slightly less than unity (and predicted to be around 0.9) in the SM. We would like to mention here that the polarization measurement in all other vector–vector modes, observed so far, (e.g., $B \to \rho K^*$ and $B \to \rho\rho$) are in accordance with the SM expectations.

Speculation on the existence of new physics in $B \to \phi K^*$ can be found in [7], in the context of two scenarios beyond the standard model (namely, R-parity violating supersymmetry and the vector-like down quark model). The issue of longitudinal polarization problem in the $B \to \phi K^*$ process and its implications were outlined in the review talk [8]. Recently, there have been a lot of works on this issue [9–15] both in and beyond the standard model. To be more specific, what we essentially need is a destructive longitudinal component, or an enhancement in the transverse component or the occurrence of both of them coherently to account for the low longitudinal polarization (or large transverse polarization) observed in $B \to \phi K^*$. But, at this stage, because of our lack of complete understanding of quark–hadron dynamics, it will be very hard to pinpoint the exact nature of it. Nevertheless, at least, it will be immensely rewarding to see if one can really afford to have a similar behavior in some of the popular scenarios beyond the SM. Therefore, in this paper we would like to analyze the decay mode $B_d^0 \to \phi K^*$ in one of the most popular scenarios beyond the SM, i.e., in the minimal supersymmetric standard model with the mass insertion approximation [16, 17] and to see whether the observed polarization anomaly can be accounted for in this model or not. It should be noted here that the contributions arising from a gluonic dipole operator with a squark–gluino loop are enhanced by a factor of $(m_{\tilde g}/m_b)$ compared to the SM contributions due to the chirality flip from the internal gluino propagator in the loop and it will interfere destructively with the longitudinal component of SM amplitude for the RL mass insertion (which is exactly what one needs, as noted earlier). Therefore, one would naively expect that the polarization anomaly in the $B \to \phi K^*$ mode can possibly be explained by the minimal supersymmetric standard model with RL mass insertion.

This paper in outline is as follows. In the next section we present the basic formalism for the $B \to V_1 V_2$ decay mode. In Sect. 3, we calculate the SM contribution in the QCD factorization approach for the sake of completeness. Section 4 contains the new physics contribution to account for the lower longitudinal polarization and in Sect. 5, we present our conclusions.

2 Polarization fractions and triple product asymmetries in $B \to V_1 V_2$ decay

The most general covariant amplitude for the decay mode $\bar B_d^0 \to V_1 V_2$ can be described as [19]

$$A(\bar B_d^0(p) \to V_1(p_1, \varepsilon_1) V_2(p_2, \varepsilon_2)) \tag{3}$$
$$= \varepsilon_{1\mu}^* \varepsilon_{2\nu}^* \left[a g^{\mu\nu} + \frac{b}{m_1 m_2} p_1^\mu p_2^\nu + \frac{\mathrm{i}c}{m_1 m_2} \epsilon^{\mu\nu\alpha\beta} p_{1\alpha} p_{2\beta} \right],$$

where p is the B-meson momentum and m_i, p_i and ε_i ($i = 1, 2$) denote the masses, momenta and polarization vectors of the outgoing vector mesons.

However, it is customary to express the angular distribution of $\bar B_d^0 \to V_1 V_2$, with each vector meson subsequently decaying into two particles, in terms of the helicity amplitudes usually defined as

$$H_\lambda = \langle V_1(\lambda) V_2(\lambda) | \mathcal{H}_{\rm eff} | \bar B_d^0 \rangle, \tag{4}$$

for $\lambda = 0, \pm 1$.

The relationship between the helicity amplitudes and the invariant amplitudes a, b, and c are given by

$$H_{\pm 1} = a \pm c\sqrt{x^2 - 1}, \quad H_0 = -ax - b(x^2 - 1), \tag{5}$$

where $x = (p_1 \cdot p_2)/m_1 m_2 = (m_B^2 - m_1^2 - m_2^2)/(2 m_1 m_2)$. The corresponding decay rate using the helicity basis amplitudes can be given by

$$\Gamma = \frac{p_c}{8\pi m_B^2} \left(|H_{+1}|^2 + |H_{-1}|^2 + |H_0|^2 \right), \tag{6}$$

where p_c is the magnitude of the CM momentum of the outgoing vector particles. It is also convenient to express the relative decay rates with longitudinal and transverse polarizations as

$$f_{\rm L} = \frac{\Gamma_{\rm L}}{\Gamma} = \frac{|H_0|^2}{|H_{+1}|^2+|H_{-1}|^2+|H_0|^2}\,, \qquad f_{\rm T} = \frac{\Gamma_{\rm T}}{\Gamma} = \frac{|H_{+1}|^2+|H_{-1}|^2}{|H_{+1}|^2+|H_{-1}|^2+|H_0|^2}\,. \tag{7}$$

The helicity amplitudes $\bar{H}_\lambda$ for the decay $B_d^0 \to \bar{V}_1\bar{V}_2$, where $\bar{V}_1$ and $\bar{V}_2$ are the antiparticles of V_1 and V_2 respectively, have the same decomposition as (3) with $a \to \bar{a}$, $b \to \bar{b}$ and $c \to -\bar{c}$. $\bar{a}$, $\bar{b}$, and $\bar{c}$ can be obtained from a, b and c by changing the sign of their weak phases.

To take the advantage of more easily extracting the CP-odd and CP-even components, the angular distribution is often written in the transversity basis. The amplitudes in transversity and helicity bases are related through the relations

$$A_\perp = \frac{H_{+1}-H_{-1}}{\sqrt{2}}\,, \qquad A_\parallel = \frac{H_{+1}+H_{-1}}{\sqrt{2}}\,, \qquad A_0 = H_0\,. \tag{8}$$

In the transversity basis the longitudinal and the CP-odd polarizations are given by

$$f_{\rm L} = \frac{|A_0|^2}{|A_\perp|^2+|A_\parallel|^2+|A_0|^2}\,, \qquad f_\perp = \frac{|A_\perp|^2}{|A_\perp|^2+|A_\parallel|^2+|A_0|^2}\,. \tag{9}$$

The triple product asymmetries (TPAs) in $\bar{B}_d^0 \to V_1V_2$ decays are defined as [20]

$$\mathcal{A}_{\rm T}^0 = \frac{1}{2}\left[\frac{{\rm Im}\,(A_\perp A_0^*)}{\sum_\lambda |A_\lambda|^2} + \frac{{\rm Im}\left(\bar{A}_\perp \bar{A}_0^*\right)}{\sum_\lambda |\bar{A}_\lambda|^2}\right]\,, \qquad \mathcal{A}_{\rm T}^\parallel = \frac{1}{2}\left[\frac{{\rm Im}\left(A_\perp A_\parallel^*\right)}{\sum_\lambda |A_\lambda|^2} + \frac{{\rm Im}\left(\bar{A}_\perp \bar{A}_\parallel^*\right)}{\sum_\lambda |\bar{A}_\lambda|^2}\right]\,, \tag{10}$$

where $\lambda = 0, \|, \perp$.

3 Standard model contribution

In the SM, the decay process $\bar{B}_d^0 \to \phi\bar{K}^{*0}$ receives a contribution from the quark level transition $b \to s\bar{s}s$, which is induced by the QCD, electroweak and magnetic penguins. The effective Hamiltonian describing the decay $b \to s\bar{s}s$ [21,22] is given by

$$\mathcal{H}_{\rm eff} = \frac{G_{\rm F}}{\sqrt{2}} V_{qb}V_{qs}^* \left[\sum_{j=3}^{10} C_jO_j + C_gO_g\right]\,, \tag{11}$$

where $q = u,\ c$. $O_3,\ldots,O_6$ and $O_7,\ldots,O_{10}$ are the standard model QCD and electroweak penguin operators respectively, and $O_{\rm g}$ is the gluonic magnetic penguin operator. The values of the Wilson coefficients at the scale $\mu \approx m_b$ in the NDR scheme are given in [23] as

$$\begin{aligned} &C_1 = 1.082\,, && C_2 = -0.185\,, && C_3 = 0.014\,,\\ &C_4 = -0.035 && C_5 = 0.009\,, && C_6 = -0.041\,,\\ &C_7 = -0.002\alpha\,, && C_8 = 0.054\alpha\,, && C_9 = -1.292\alpha\,,\\ &C_{10} = 0.263\alpha\,, && C_{\rm g} = -0.143\,. \end{aligned} \tag{12}$$

We use the QCD factorization approach [21,22] to evaluate the hadronic matrix elements, which allows us to compute the non-factorizable corrections in the heavy quark limit. Naive factorization is recovered in the heavy quark limit and to the zeroth order of the QCD corrections. The decay mode $B \to \phi K^*$ has been analyzed in [15,18,24] using the QCD factorization approach. We will first briefly discuss the essential differences between these three approaches. It has been shown in [24] that the magnetic dipole penguin will contribute to all the three helicity amplitudes with almost the same order. Later, this result has been changed in [18] where they have shown that the magnetic penguin will contribute only to the longitudinal polarization amplitude (H_0). However, very recently again there has been a correction in [15], namely that the positive helicity amplitude (H_{+1}) also receives small but non-zero contributions from the magnetic dipole operator. It should be noted here that the contribution of the magnetic dipole operator to the tranverse amplitudes are also found to be small but non-zero in the pQCD approach by Li and Mishima [14].

Here, we will use the results of the QCD factorization method as obtained in [15]. In this approach the helicity amplitudes are given by

$$\begin{aligned} H_0 &= -\frac{G_{\rm F}}{\sqrt{2}} V_{tb}V_{ts}^*\,\frac{\tilde{a}^0 f_\phi}{2m_{K^*}}\\ &\times\Bigg[(m_B^2 - m_{K^*}^2 - m_\phi^2)(m_B+m_{K^*})A_1^{BK^*}(m_\phi^2)\\ &\quad - \frac{4m_B^2p_c^2}{m_B+m_{K^*}}A_2^{BK^*}(m_\phi^2)\Bigg]\,,\\ H_{\pm1} &= -\frac{G_{\rm F}}{\sqrt{2}} V_{tb}V_{ts}^*\;\tilde{a}^\pm m_\phi f_\phi\\ &\times\left[(m_B+m_{K^*})A_1^{BK^*}(m_\phi^2) \mp \frac{2m_Bp_c}{m_B+m_{K^*}}V^{BK^*}(m_\phi^2)\right]\,, \end{aligned} \tag{13}$$

where $\tilde{a}^h = a_3^h + a_4^h + a_5^h - \frac{1}{2}(a_7^h + a_9^h + a_{10}^h)$ with $h = 0, \pm1$. $A_{1,2}^{BK^*}(q^2)$ and $V^{BK^*}(q^2)$ are the form factors describing the $B \to K^*$ transitions [25] evaluated at $q^2 = m_\phi^2$. The expressions for the effective parameters a_i^h appearing in the helicity amplitudes (13) are given by [15]

$$\begin{aligned} a_3^h &= C_3 + \frac{C_4}{N} + \frac{\alpha_{\rm s}}{4\pi}\frac{C_F}{N}C_4\Big(f_{\rm I}^h(1) + f_{\rm II}^h(1)\Big)\,,\\ a_4^h &= C_4 + \frac{C_3}{N} \end{aligned}$$

$$
\begin{aligned}
&+\frac{\alpha_s}{4\pi}\frac{C_F}{N}\left(C_3\left[f_{\rm I}^h(1)+f_{\rm II}^h(1)+G^h(s_s)+G^h(s_b)\right]\right.\\
&-C_1\left(\frac{\lambda_u}{\lambda_t}G^h(s_u)+\frac{\lambda_c}{\lambda_t}G^h(s_c)\right)\\
&+(C_4+C_6)\sum_{i=u}^{b}\left(G^h(s_i)-\frac{2}{3}\right)\\
&+\frac{3}{2}(C_8+C_{10})\sum_{i=u}^{b}e_i\left(G^h(s_i)-\frac{2}{3}\right)\\
&\left.+\frac{3}{2}C_9\left[e_sG^h(s_s)+e_bG^h(s_b)\right]+C_{\rm g}G_{\rm g}^h\right),\\
a_5^h&=C_5+\frac{C_6}{N}-\frac{\alpha_s}{4\pi}\frac{C_F}{N}C_6\Big[f_{\rm I}^h(-1)+f_{\rm II}^h(-1)\Big],\\
a_7^h&=C_7+\frac{C_8}{N}-\frac{\alpha_s}{4\pi}\frac{C_F}{N}C_8\Big[f_{\rm I}^h(-1)+f_{\rm II}^h(-1)\Big]\\
&-\frac{\alpha}{9\pi}NC_e^h,\\
a_9^h&=C_9+\frac{C_{10}}{N}+\frac{\alpha_s}{4\pi}\frac{C_F}{N}C_{10}\Big[f_{\rm I}^h(1)+f_{\rm II}^h(1)\Big]\\
&-\frac{\alpha}{9\pi}NC_e^h,\\
a_{10}^h&=C_{10}+\frac{C_9}{N}+\frac{\alpha_s}{4\pi}\frac{C_F}{N}C_9\Big[f_{\rm I}^h(1)+f_{\rm II}^h(1)\Big]\\
&-\frac{\alpha}{9\pi}C_e^h,
\end{aligned} \tag{14}
$$

where $\lambda_q=V_{qb}V_{qs}^*$, $C_F=(N^2-1)/2N$ and $s_i=m_i^2/m_b^2$. The QCD penguin loop functions $G^h(s)$ are given by

$$
\begin{aligned}
G^0(s)&=\frac{2}{3}-\frac{4}{3}\ln\frac{\mu}{m_b}+4\int_0^1{\rm d}x\ \Phi_\parallel^V(x)\ g(x,s),\\
G^{\pm1}(s)&=\frac{2}{3}-\frac{2}{3}\ln\frac{\mu}{m_b}\\
&+2\int_0^1{\rm d}x\ \left(g_\perp^{(v)\phi}(x)\pm\frac{1}{4}\frac{{\rm d}g_\perp^{(a)\phi}(x)}{{\rm d}x}\right)g(x,s),
\end{aligned} \tag{15}
$$

with the function $g(x,s)$ defined as

$$
g(x,s)=\int_0^1{\rm d}u\ u(1-u)\ln\left[s-u(1-u)(1-x)-{\rm i}\epsilon\right]. \tag{16}
$$

The EW penguin type diagrams induced by the operators O_1 and O_2 are

$$
C_e^h=\left(\frac{\lambda_u}{\lambda_t}G^h(s_u)+\frac{\lambda_c}{\lambda_t}G^h(s_c)\right)\left(C_2+\frac{C_1}{N}\right). \tag{17}
$$

The gluonic dipole operator O_g gives a tree level contribution as

$$
\begin{aligned}
G_g^0&=-2\int_0^1{\rm d}x\ \frac{\Phi_\parallel^\phi(x)}{1-x},\\
G_g^+&=-\int_0^1{\rm d}x\left(g_\perp^{(v)\phi}(x)+\frac{1}{4}\frac{{\rm d}g_\perp^{(a)\phi}(x)}{{\rm d}x}\right)\frac{1}{1-x},\\
G_g^-&=0.
\end{aligned} \tag{18}
$$

The vertex correction factors $f_{\rm I}^h$ are given by

$$
\begin{aligned}
f_{\rm I}^0(a)&=-12\ln\frac{\mu}{m_b}-18+6(1-a)\\
&+\int_0^1{\rm d}x\ \Phi_\parallel^\phi(x)\left(3\frac{1-2x}{1-x}\ln x-3{\rm i}\pi\right),\\
f_{\rm I}^{\pm1}(a)&=-12\ln\frac{\mu}{m_b}-18+6(1-a)\\
&+\int_0^1{\rm d}x\ \left(g_\perp^{(v)\phi}(x)\pm\frac{a}{4}\frac{{\rm d}g_\perp^{(a)\phi}(x)}{{\rm d}x}\right)\\
&\times\left(3\frac{1-2x}{1-x}\ln x-3{\rm i}\pi\right).
\end{aligned} \tag{19}
$$

The hard spectator interaction $f_{\rm II}^h$ arising from the hard spectator interaction with a hard gluon exchange between the vector meson and the spectator quark of the B-meson is given by

$$
\begin{aligned}
f_{\rm II}^0(a)&=\frac{4\pi^2}{N}\frac{{\rm i}f_Bf_{K^*}f_\phi}{h_0}\\
&\times\int_0^1{\rm d}\rho\frac{\Phi_1^B(\rho)}{\rho}\int_0^1{\rm d}v\frac{\Phi_\parallel^{K^*}(v)}{\bar v}\int_0^1{\rm d}u\frac{\Phi_\parallel^\phi(u)}{u},\\
f_{\rm II}^{\pm1}(a)&=-\frac{4\pi^2}{N}\frac{2{\rm i}f_Bf_{K^*}^\perp f_\phi m_\phi}{m_Bh_{\pm1}}(1\mp1)\\
&\times\int_0^1{\rm d}\rho\frac{\Phi_1^B(\rho)}{\rho}\int_0^1{\rm d}v\frac{\Phi_\perp^{K^*}(v)}{\bar v^2}\\
&\times\int_0^1{\rm d}u\left(g_\perp^{(v)\phi}(u)-\frac{a}{4}\frac{{\rm d}g_\perp^{(a)\phi}(u)}{{\rm d}u}\right)\\
&+\frac{4\pi^2}{N}\frac{2{\rm i}f_Bf_{K^*}f_\phi m_{K^*}m_\phi}{m_B^2h_{\pm1}}\\
&\times\int_0^1{\rm d}\rho\frac{\Phi_1^B(\rho)}{\rho}\int_0^1{\rm d}v{\rm d}u\left(g_\perp^{(v)K^*}(v)\pm\frac{1}{4}\frac{{\rm d}g_\perp^{(a)K^*}(v)}{{\rm d}v}\right)\\
&\times\left(g_\perp^{(v)\phi}(u)\pm\frac{a}{4}\frac{{\rm d}g_\perp^{(a)\phi}(u)}{{\rm d}u}\right)\frac{u+\bar v}{u\bar v^2},
\end{aligned} \tag{20}
$$

with $\bar v=1-v$ and

$$
\begin{aligned}
h_0&=\frac{{\rm i}f_\phi}{2m_{K^*}}\Big[(m_B^2-m_{K^*}^2-m_\phi^2)(m_B+m_{K^*})A_1^{BK^*}(m_\phi^2)\\
&-\frac{4m_B^2p_c^2}{m_B+m_{K^*}}A_2^{BK^*}(m_\phi^2)\Big],\\
h_{\pm1}&={\rm i}f_\phi m_\phi\Big[(m_B+m_{K^*})A_1^{BK^*}(m_\phi^2)
\end{aligned}
$$

$$\mp \frac{2m_B p_c}{m_B + m_{K^*}} V^{BK^*}(m_\phi^2) \Bigg] . \quad (21)$$

The asymptotic form of the leading twist ($\Phi_\parallel^V(x)$, $\Phi_\perp^V(x)$) and twist-3 ($g_\perp^{(v)}(x)$, $g_\perp^{(a)}(x)$) light-cone distribution amplitudes are defined as

$$\Phi_\parallel^V(x) = \Phi_\perp^V(x) = g_\perp^{(a)}(x) = 6x(1-x) ,$$
$$g_\perp^{(v)}(x) = \frac{3}{4}[1 + (2x-1)^2] . \quad (22)$$

The light-cone projector for the B-meson in the heavy quark limit can be expressed as [21]

$$\mathcal{M}^B = -\frac{\mathrm{i} f_B m_B}{4} \left[(1+ \not{v})\gamma_5 \left\{\Phi_1^B(\xi) + \not{n}_- \Phi_2^B(\xi)\right\}\right] , \quad (23)$$

where ξ is the momentum fraction of the spectator quark in the B-meson, $v = (1,0,0,0)$, $n_- = (1,0,0,-1)$ is the light-cone vector. The normalization conditions are given by

$$\int_0^1 \mathrm{d}\xi \ \Phi_1^B(\xi) = 1 , \quad \int_0^1 \mathrm{d}\xi \ \Phi_2^B(\xi) = 0 . \quad (24)$$

For our numerical evaluation we use

$$\int_0^1 \mathrm{d}\xi \frac{\Phi_1^B(\xi)}{\xi} = \frac{m_B}{\lambda_B} , \quad (25)$$

with $\lambda_B = 0.46\,\mathrm{GeV}$, which parametrizes our ignorance of the B-meson distribution amplitudes.

It should be noted that the presence of logarithmic and linear infrared divergences in $f_{\mathrm{II}}^{\pm 1}$ implies that the spectator interaction is dominated by the soft gluon exchanges in the final states. To regulate these divergences, a cutoff parameter of order $\Lambda_{\mathrm{QCD}}/m_b$, with $\Lambda_{\mathrm{QCD}} = 0.5\,\mathrm{GeV}$ has been used.

For our numerical analysis, we use the following input parameters. The quark masses appearing in the penguin diagrams are pole masses and we have used the following values (in GeV): $m_u = m_d = m_s = 0$, $m_c = 1.4$ and $m_b = 4.8$. The decay constants used are (in GeV) $f_B = 0.161$, $f_{K^*} = 0.217$, $f_\phi = 0.231$ and $f_{K^*}^\perp = 0.156$. The form factors are evaluated in the light-cone sum rule analysis [26] where the q^2 dependence is given by

$$F(q^2) = F(0) \exp[c_1(q^2/m_B^2) + c_2(q^2/m_B^2)^2] , \quad (26)$$

with the parameters as given in Table 2. The particle masses and lifetime of B_d^0-meson have been taken from [27]. For the CKM matrix elements, we have used [27]

$$|V_{cb}| = 0.0413 \pm 0.0015 , \quad \bar{\rho} = 0.20 \pm 0.09 ,$$

Table 2. The parameters of the form factors describing $B \to K^*$ transitions

	$A_1(0)$	$A_2(0)$	$V(0)$
$F(0)$	0.294	0.246	0.399
c_1	0.656	1.237	1.537
c_2	0.456	0.822	1.123

$$\bar{\rho} = 0.33 \pm 0.05 .$$

With these input parameters, we obtain the branching ratio in the SM as

$$\mathcal{B}(B_d^0 \to \phi K^{*0}) = (6.34 \pm 0.46) \times 10^{-6} , \quad (27)$$

and the longitudinal and the CP-odd polarizations as

$$f_{\mathrm{L}} = 0.89 , \quad f_\perp = 0.05 . \quad (28)$$

The triple product asymmetries $\mathcal{A}_{\mathrm{T}}^{(0,\parallel)}$ (10) are found to be identically zero.

4 New physics contributions

We now consider the contribution arising from NP. In general the effective $\Delta B = 1$, NP Hamiltonian relevant for the $b \to s\bar{s}s$ transition is given by

$$\mathcal{H}_{\mathrm{eff}}^{\mathrm{NP}} \propto \left[\sum_i (C_i^{\mathrm{NP}} O_i + \tilde{C}_i^{\mathrm{NP}} \tilde{O}_i) + C_g O_g + \tilde{C}_g \tilde{O}_g\right] , \quad (29)$$

where O_i (O_g), are the standard model like QCD (magnetic) penguin operators with current structure $(\bar{s}b)_{V-A}(\bar{s}s)_{V\pm A}$ and C_i^{NP}, C_g^{NP} are the new Wilson coefficients. The operators $\tilde{O}_i$ ($\tilde{O}_g$) are obtained from O_i (O_g) by exchanging $L \leftrightarrow R$. As discussed in [9] the NP contributions to the different helicity amplitudes are given by

$$A^{\mathrm{NP}}(\bar{B}_d^0 \to \phi K^{*0})_{0,\parallel} \propto C_i^{\mathrm{NP}} - \tilde{C}_i^{\mathrm{NP}} ,$$
$$A^{\mathrm{NP}}(\bar{B}_d^0 \to \phi K^{*0})_{\perp} \propto C_i^{\mathrm{NP}} + \tilde{C}_i^{\mathrm{NP}} . \quad (30)$$

Thus, in the presence of new physics, the different amplitudes can be given by

$$A_{0,\parallel} = A_{0,\parallel}^{\mathrm{SM}} + A_{0,\parallel}^{\mathrm{NP}} = A_{0,\parallel}^{\mathrm{SM}} \left[1 + \mathrm{e}^{\mathrm{i}\phi_N}(r_{0,\parallel} - \tilde{r}_{0,\parallel})\right] ,$$
$$A_\perp = A_\perp^{\mathrm{SM}} + A_\perp^{\mathrm{NP}} = A_\perp^{\mathrm{SM}} \left[1 + \mathrm{e}^{\mathrm{i}\phi_N}(r_\perp + \tilde{r}_\perp)\right] , \quad (31)$$

where r_λ, with $(\lambda = 0, \parallel, \perp)$ are the ratio of NP (arising from $C_i O_i$ and $C_g O_g$ part of the Hamiltonian) to SM amplitudes, $\tilde{r}_\lambda$ are the corresponding values arising from the $\tilde{C}_i \tilde{O}_i$ and $\tilde{C}_g \tilde{O}_g$ part. ϕ_N is the relative weak phase between the SM and NP amplitudes. For simplicity, we have assumed a common weak phase for the C and $\tilde{C}$ contributions and a zero strong phase between the SM and the NP amplitudes.

Thus the branching ratio is given by

$$\mathcal{B}(B_d^0 \to \phi K^{*0}) \quad (32)$$
$$= \mathcal{B}^{\mathrm{SM}} \left[1 + \frac{\sum_\lambda R_\lambda^2 |A_\lambda^{\mathrm{SM}}|^2}{\sum_\lambda |A_\lambda^{\mathrm{SM}}|^2} + 2\frac{\sum_\lambda R_\lambda |A_\lambda^{\mathrm{SM}}|^2}{\sum_\lambda |A_\lambda^{\mathrm{SM}}|^2} \cos\phi_N\right] ,$$

where $R_{\parallel,0} = r_{\parallel,0} - \tilde{r}_{\parallel,0}$ and $R_\perp = r_\perp + \tilde{r}_\perp$ and $\mathcal{B}^{\mathrm{SM}}$ denotes the SM branching ratio. The longitudinal and the CP-odd polarizations now read as

$$f_{\mathrm{L}} = \frac{|A_0^{\mathrm{SM}}|^2 \left[1 + R_0^2 + 2R_0 \cos\phi_N\right]}{\sum_\lambda |A_\lambda^{\mathrm{SM}}|^2 \left[1 + R_\lambda^2 + 2R_\lambda \cos\phi_N\right]} ,$$

$$f_\perp = \frac{|A_\perp^{\rm SM}|^2 \left[1 + R_\perp^2 + 2R_\perp \cos\phi_N\right]}{\sum_\lambda |A_\lambda^{\rm SM}|^2 \left[1 + R_\lambda^2 + 2R_\lambda \cos\phi_N\right]} \,. \quad (33)$$

Furthermore, in the presence of NP, the triple product asymmetries (10) are given by

$$\mathcal{A}_{\rm T}^0 = \frac{2(R_\perp - R_0)\sin\phi_N}{\sum_\lambda |A_\lambda^{\rm SM}|^2 \left[1 + R_\lambda^2 + 2R_\lambda \cos\phi_N\right]} \,,$$
$$\mathcal{A}_{\rm T}^{\|} = \frac{2(R_\perp - R_\|)\sin\phi_N}{\sum_\lambda |A_\lambda^{\rm SM}|^2 \left[1 + R_\lambda^2 + 2R_\lambda \cos\phi_N\right]} \,. \quad (34)$$

We now analyze the decay process $B_d^0 \to \phi K^{*0}$ in the minimal supersymmetric standard model (MSSM) with mass insertion approximation. This decay mode receives supersymmetric (SUSY) contributions mainly from penguin and box diagrams containing gluino–squark, chargino–squark and charged Higgs–top loops. Here, we consider only the gluino contributions, because the chargino and charged Higgs loops are expected to be suppressed by the small electroweak gauge couplings. However, the gluino mediated FCNC contributions are of the order of the strong interaction strength, which may exceed the existing limits. Therefore, it is customary to rotate the effects, so that the FCNC effects occur in the squark propagators rather than in the couplings and to parameterize them in terms of dimensionless parameters. Here we work in the usual mass insertion approximation [16, 17], where the flavor mixings $i \to j$ in the down-type squarks associated with $\tilde{q}_B$ and $\tilde{q}_A$ are parametrized by $(\delta_{AB}^d)_{ij}$, with $A, B = L, R$ and i, j as the generation indices. More explicitly $(\delta_{LL}^d)_{ij} = (V_L^{d\dagger} M_{\tilde{d}}^2 V_L^d)_{ij}/m_{\tilde{q}}^2$, where $M_{\tilde{d}}^2$ is the squared down squark mass matrix and $m_{\tilde{q}}$ is the average squark mass.

V_d is the matrix which diagonalizes the down-type quark mass matrix.

Thus, the new effective $\Delta B = 1$ Hamiltonian relevant for the $B_d^0 \to \phi K^*$ process arising from new penguin/box diagrams with gluino–squark in the loops is given by

$$\mathcal{H}_{\rm eff}^{\rm SUSY} = -\frac{G_{\rm F}}{\sqrt{2}} V_{tb} V_{ts}^* \times \left[\sum_{i=3}^{6}\left(C_i^{\rm NP} O_i + \tilde{C}_i^{\rm NP} \tilde{O}_i\right) + C_g^{\rm NP} O_g + \tilde{C}_g^{\rm NP} \tilde{O}_g\right] , \quad (35)$$

where O_i (O_g) are the QCD (magnetic) penguin operators and the $C_i^{\rm NP}$ (C_g) are the new Wilson coefficients. The operators $\tilde{O}_i$ are obtained from O_i by exchanging $L \leftrightarrow R$.

To evaluate the amplitude in the MSSM, we have to first determine the Wilson coefficients at the b quark mass scale. At the leading order, in the mass insertion approximation, the new Wilson coefficients corresponding to each of the operator at the scale $\mu \sim \tilde{m} \sim M_W$ are given by [17, 28]

$$C_3^{\rm NP} \simeq -\frac{\sqrt{2}\alpha_{\rm s}^2}{4G_{\rm F} V_{tb} V_{ts}^* m_{\tilde{q}}^2} \left(\delta_{LL}^d\right)_{23} \times \left[-\frac{1}{9}B_1(x) - \frac{5}{9}B_2(x) - \frac{1}{18}P_1(x) - \frac{1}{2}P_2(x)\right] ,$$

$$C_4^{\rm NP} \simeq -\frac{\sqrt{2}\alpha_{\rm s}^2}{4G_{\rm F} V_{tb} V_{ts}^* m_{\tilde{q}}^2} \left(\delta_{LL}^d\right)_{23} \times \left[-\frac{7}{3}B_1(x) + \frac{1}{3}B_2(x) + \frac{1}{6}P_1(x) + \frac{3}{2}P_2(x)\right] ,$$

$$C_5^{\rm NP} \simeq -\frac{\sqrt{2}\alpha_{\rm s}^2}{4G_{\rm F} V_{tb} V_{ts}^* m_{\tilde{q}}^2} \left(\delta_{LL}^d\right)_{23} \times \left[\frac{10}{9}B_1(x) + \frac{1}{18}B_2(x) - \frac{1}{18}P_1(x) - \frac{1}{2}P_2(x)\right] ,$$

$$C_6^{\rm NP} \simeq -\frac{\sqrt{2}\alpha_{\rm s}^2}{4G_{\rm F} V_{tb} V_{ts}^* m_{\tilde{q}}^2} \left(\delta_{LL}^d\right)_{23} \times \left[-\frac{2}{3}B_1(x) + \frac{7}{6}B_2(x) + \frac{1}{6}P_1(x) + \frac{3}{2}P_2(x)\right]$$

$$C_{\rm g}^{\rm NP} \simeq -\frac{2\sqrt{2}\pi\alpha_{\rm s}}{2G_{\rm F} V_{tb} V_{ts}^* m_{\tilde{q}}^2} \left[\left(\delta_{LL}^d\right)_{23}\left(\frac{3}{2}M_3(x) - \frac{1}{6}M_4(x)\right) + \left(\delta_{LR}^d\right)_{23}\left(\frac{m_{\tilde{g}}}{m_b}\right)\frac{1}{6}\left(4B_1(x) - \frac{9}{x}B_2(x)\right)\right] , \quad (36)$$

where $x = m_{\tilde{g}}^2/m_{\tilde{q}}^2$. The loop functions appearing in these expressions can be found in [17]. The corresponding $\tilde{C}_i^{\rm NP}$ are obtained from $C_i^{\rm NP}$ by interchanging $L \leftrightarrow R$. It should be noted that the $(\delta_{LR}^d)_{23}$ contribution is enhanced by $(m_{\tilde{g}}/m_b)$ compared to that of the SM and the LL insertion due to the chirality flip from the internal gluino propagator in the loop. Therefore, the magnetic dipole operators in supersymmetric model are found to contribute significantly.

The Wilson coefficients at low energy, $C_i^{\rm NP}(\mu \sim m_b)$, can be obtained from $C_i^{\rm NP}(M_W)$ by using the renormalization group (RG) equation, as discussed in [23], as

$$\boldsymbol{C}(\mu) = \boldsymbol{U}_5(\mu, M_W)\boldsymbol{C}(M_W) \,, \quad (37)$$

where $\boldsymbol{C}$ is the 6×1 column vector of the Wilson coefficients and $\boldsymbol{U}_5(\mu, M_W)$ is the five-flavor 6×6 evolution matrix. In the next-to-leading order (NLO), $\boldsymbol{U}_5(\mu, M_W)$ is given by

$$\boldsymbol{U}_5(\mu, M_W) = \left(1 + \frac{\alpha_{\rm s}(\mu)}{4\pi}\boldsymbol{J}\right)\boldsymbol{U}_5^{(0)}(\mu, M_W)\left(1 - \frac{\alpha_{\rm s}(M_W)}{4\pi}\boldsymbol{J}\right) , \quad (38)$$

where $\boldsymbol{U}_5^{(0)}(\mu, M_W)$ is the leading order (LO) evolution matrix and $\boldsymbol{J}$ denotes the NLO corrections to the evolution. The explicit forms of $\boldsymbol{U}_5(\mu, M_W)$ and $\boldsymbol{J}$ are given in [23].

Since the O_g contribution to the matrix element is $\alpha_{\rm s}$ order suppressed, we consider only leading order RG effects for the coefficient $C_g^{\rm NP}$, which is given by [28]

$$C_g^{\rm NP}(m_b) \simeq -0.15 + 0.70\, C_{\rm g}^{\rm NP}(M_W) \,. \quad (39)$$

For the numerical analysis, we fix the SUSY parameter as $m_{\tilde{q}} = m_{\tilde{g}} = 500\,{\rm GeV}$, $\alpha_{\rm s}(M_W) = 0.119$. The absolute values of the mass insertion parameters $(\delta_{AB}^d)_{23}$, with $A, B = (L, R)$, are constrained by the experimental value of the $B \to X_s\gamma$ decay [17]. These constraints are very

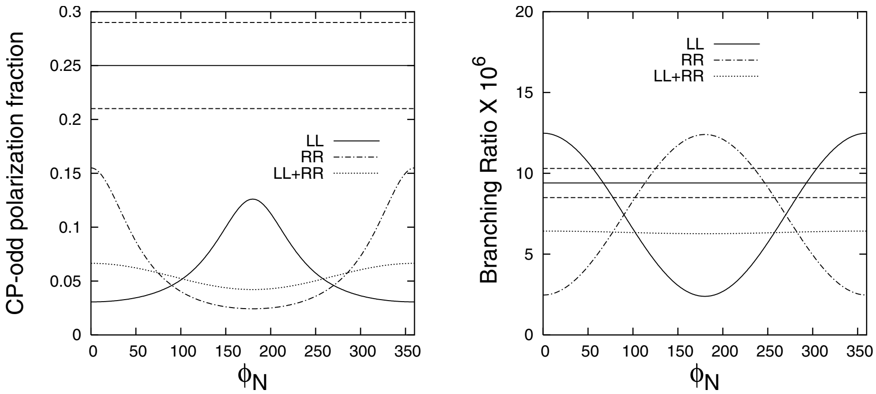

Fig. 1. The CP-odd polarization fraction ($f_\perp$) and the branching ratio of the process $B_d^0 \to \phi K^{*0}$ with LL and RR mass insertions versus the weak phase ϕ_N (in degree). The horizontal solid lines represent the experimental central value and the dashed lines represent the 1σ range

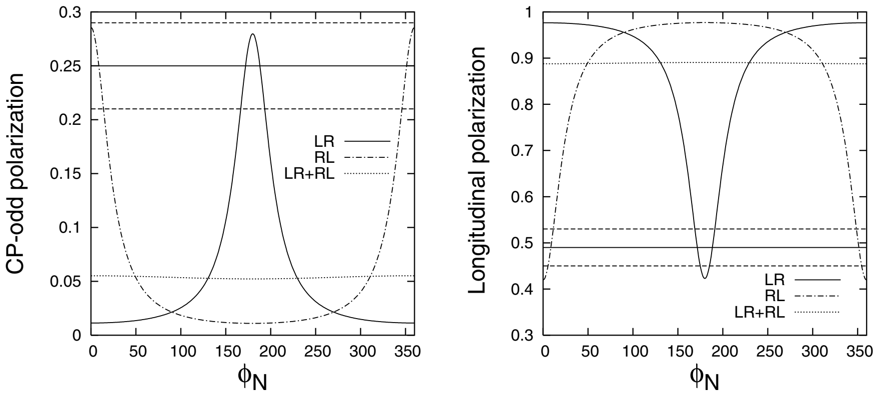

Fig. 2. The CP-odd and the longitudinal polarization fractions of the process $B_d^0 \to \phi K^{*0}$ with LR and RL mass insertions

weak for LL and RR mass insertions and the existing limits come only from their definitions $|(\delta^d_{LL,RR})_{23}| < 1$. The LR and RL mass insertions are more constrained and for instance with $m_{\tilde g} \simeq m_{\tilde q} \simeq 500\,\text{GeV}$, we have $|(\delta^d_{LR,RL})_{23}| \leq 1.6 \times 10^{-2}$. In our analysis, we will use the above bounds on the mass insertion parameters, i.e.,

$$|(\delta^d_{LL,RR})_{23}| < 1 \quad \text{and} \quad |(\delta^d_{LR,RL})_{23}| \leq 1.6 \times 10^{-2}\,. \quad (40)$$

Now substituting the values of the RG evolved Wilson coefficients $C_i^{\mathrm{NP}}(m_b)$ in (14) we obtain the corresponding a_i^h's and hence with (8), (13) and (31) the amplitudes. Assuming that all the mass insertion parameters $(\delta^d_{AB})_{23}$ have a common weak phase, we obtain the new physics parameters arising from the LL (LR) and RR (RL) mass insertions as

$$\begin{aligned}
&(r_0)_{LL} = (\tilde r_0)_{RR} < 0.44\,, \quad (r_0)_{LR} = (\tilde r_0)_{RL} \leq 1.3\,,\\
&(r_\parallel)_{LL} = (\tilde r_\parallel)_{RR} < 6.2 \times 10^{-2}\,,\\
&(r_\parallel)_{LR} = (\tilde r_\parallel)_{RL} \leq 7.0 \times 10^{-2}\,, \qquad (41)\\
&(r_\perp)_{LL} = (\tilde r_\perp)_{RR} < 6.0 \times 10^{-2}\,,\\
&(r_\perp)_{LR} = (\tilde r_\perp)_{RL} \leq 7.2 \times 10^{-2}\,.
\end{aligned}$$

Let us now analyze the variation of the CP-odd polarization fraction and the branching ratio in the presence of new physics. We first consider the contributions arising from the LL and RR mass insertions. As seen from (41), these contributions are quite small and it is expected that they cannot accommodate the observed large CP-odd polarization. Now using the maximum values of r_λ from (41), we plot the CP-odd polarization ($f_\perp$) (33) and the branching ratio (32) versus the weak phase ϕ_N for three different cases (LL, RR and in the presence of both the LL and the RR contributions). It is seen from Fig. 1 that, indeed the large CP-odd polarization fraction cannot be accommodated in these cases although the observed branching ratio can be accommodated with LL or RR mass insertions. Next we consider the contributions arising from LR, RL and the simultaneous presence of LR and RL mass insertions. As seen from Fig. 2, in this case the observed CP-odd polarization fraction ($f_\perp$) and the longitudinal polarization fraction (f_{L}) can be accommodated with either LR or RL mass insertions. The branching ratio can also be accommodated with these mass insertions as seen from Fig. 3.

5 Conclusions

Observation of an unexpectedly small longitudinal polarization (and large transverse polarization) in the penguin dominated $B \to \phi K^*$ mode poses a serious challenge both to theorists and experimentalists in B-physics. This has in turn ignited the desire of revealing the existence of new physics beyond the SM. While at present there is no clear indication of any NP but with the accumulation of more data, if the longitudinal polarization stabilizes around the present central value, i.e., $f_{\mathrm{L}} = 0.5$, with *reduced error bars*, then this might be the first clear evidence of new physics in the $b \to s$ penguin decay amplitudes. The polarization measurements in various vector–vector modes undertaken

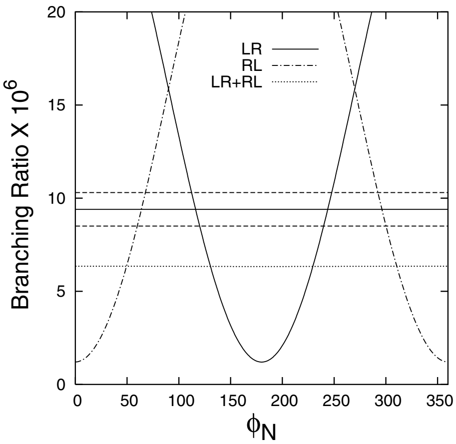

Fig. 3. The branching ratio of the process $B_d^0 \to \phi K^{*0}$ with LR and RL mass insertions

by the ongoing B-factory experiments and the experiments to be performed at BTeV and LHC-b will definitely be able to guide us to an understanding of the dynamics and possibly provide us with a meaningful answer to our quest for the existence of new physics.

We have employed supersymmetry with the mass insertion approximation and shown that the low longitudinal polarization in $B_d^0 \to \phi K^{*0}$ can be accommodated in this scenario beyond the standard model with either the LR or the RL mass insertion approximation.

Since in the B-factory data so far we have only seen some kind of deviation in the $b \to s\bar{s}s$ transition measurements it may be worthwhile to continue our effort in this direction and check carefully if we can really observe NP in this type of penguin induced transitions. If NP is present in the $b \to s\bar{s}s$ transitions and indeed if it is responsible for the observed lower longitudinal polarization then we expect to see the same effect of lower longitudinal polarization in another charmless vector–vector mode, i.e., $B_s \to \phi\phi$, which is governed by the same penguin induced $b \to s\bar{s}s$ transition. Already the branching ratio for this mode has been measured by the CDF Collaboration [29,30] and we are looking forward to the polarization measurements in this mode. This in turn, at least, will provide us with a clear picture of the charmless vector–vector transitions induced by the $b \to s\bar{s}s$ penguins and possibly revealing the existence of NP in penguin induced $b \to s\bar{s}s$ transitions. If confirmed, the polarization anomaly along with the deviation measured in $\sin 2\beta$ measurements may bring us one step further towards the establishment of NP in the penguin induced $b \to s\bar{s}s$ transitions. It is therefore urgently needed to closely examine experimentally all the possible charmless vector–vector modes to confirm or rule out the existence of new physics.

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