$B \to J/\psi K^*$ in a supersymmetric right-handed flavor mixing scenario

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A supersymmetric extension of the standard model, with maximal $\tilde{s}_R$–$\tilde{b}_R$ mixing and a new source of CP violation, contains all the necessary ingredients to account for a possible anomaly in the measured CP asymmetry in $B \to \phi K_S$ decay. In the same framework we study the decay $B \to J/\psi K^*$, paying particular attention to observables that can be extracted by performing a time-dependent angular analysis, and become nonzero because of new physics effects.

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I. INTRODUCTION

Recent results on B decay to charmonium modes, such as $B \to J/\psi K_S(K^*)$, are in good agreement with the standard model (SM). However, for the time-dependent CP asymmetry in $B \to \phi K_S$ mode, also if new results seem to show more agreement with the SM, $S_{\phi K_S} = 0.50 \pm 0.25$ (*BABAR*) [1] and $S_{\phi K_S} = 0.06 \pm 0.33$ (Belle) [2], than in the past (before ICHEP04, the world average was $S_{\phi K_S} = 0.02 \pm 0.29$ [3]), new physics (NP) effects could still play some role [4]. A deviation from the SM expectation of $S_{\phi K_S} \cong S_{\psi K_S}$ would call for large $s - b$ mixing, the existence of a new source of CP violation, and perhaps right-handed dynamics [5]. It would be important to find confirming evidence in the future.

In anticipation of a future ''super B factory'' that would allow precision measurements, we study CP violation in the vector-vector $B_d^0 \to J/\psi K^*$ mode by taking into account deviations in $S_{\phi K_S}$. We do not expect the NP effects to show up in the branching fraction since, in contrast to $B \to \phi K_S$, $B_d^0 \to J/\psi K^*$ decay is tree dominant. With special attention to observables related to the time-dependent CP asymmetry [6–8], we focus on manifestations of NP effects. In the phenomenological analysis presented in this paper, we pay particular attention to those observables which are expected to vanish in the SM [8].

II. $B \to VV$

The $B_d^0 \to J/\psi K^{(*)}$ decay is dominated by the tree-level $\bar{b} \to c\bar{c}\,\bar{s}$ process, while the CP phase in the corresponding penguin amplitude is highly suppressed. If NP contributions are present, it could manifest itself as direct CP violation effects. This can be illustrated by the full amplitude for the decay $B \to f$,

$$A(B \to f) = ae^{i\delta^a} + be^{i\phi}e^{i\delta^b}, \tag{1}$$

where the weak phases are assumed to be zero for the first amplitude and ϕ for the second, and $\delta^{a,b}$ are the respectively strong phases. For the CP conjugate decay $\bar{B} \to \bar{f}$ the amplitude is given by changing the sign of the weak phase ϕ. The direct CP asymmetry is then obtained,

$$a_{\rm dir}^{CP} = \frac{\Gamma(B \to f) - \Gamma(\bar{B} \to \bar{f})}{\Gamma(B \to f) + \Gamma(\bar{B} \to \bar{f})} = \frac{2ab\sin(\delta^a - \delta^b)\sin\phi}{a^2 + b^2 + 2ab\cos(\delta^a - \delta^b)\cos\phi}, \tag{2}$$

which does not vanish for $\phi \neq 0$ and if the strong phase difference is also not zero. Recent measurements of the direct CP asymmetries of $B^0 \to K^+\pi^-$ and $B^+ \to K^+\pi^0$, with values -0.114 ± 0.020 and $+0.049 \pm 0.040$ respectively [1], do not exclude the possibility that strong phases could be not so small, as initially thought because of the b quark being rather heavy. Whether or not effects of NP in the direct CP asymmetry will be detectable will depend on how big the strong phases are. It is therefore important that one can still seek NP effects by performing a time-dependent analysis and comparing $B_d^0(t) \to J/\psi K^*$ with $J/\psi K_S$. In fact, more information is contained in the time-dependent angular analysis of vector-vector decays such as $B_d^0(t) \to J/\psi K^*$ or ϕK^* [7,8].

For a $B \to VV$ decay, the final state can be decomposed into three helicity amplitudes $\{A_0, A_\parallel, A_\perp\}$. A_0 corresponds to both the vector mesons being polarized along their direction of motion, while $A_\parallel$ and $A_\perp$ correspond to both polarization states being transverse to their directions of motion but parallel and orthogonal to each other, respectively [9]. If, in particular, we consider the decay $B_d^0(t) \to J/\psi K^*$, analogous to Eq. (1) we have,

$$A_\lambda(B \to J/\psi K^*) = a_\lambda e^{i\delta_\lambda^a} + b_\lambda e^{i\phi}e^{i\delta_\lambda^b}, \tag{3}$$

$$\bar{A}_\lambda(\bar{B} \to J/\psi \bar{K}^*) = a_\lambda e^{i\delta_\lambda^a} + b_\lambda e^{-i\phi}e^{i\delta_\lambda^b}, \tag{4}$$

where a_λ and b_λ are the SM and NP amplitudes and $\delta_\lambda^{a,b}$ their respective strong phases, for each helicity component. The full decay amplitude becomes

$$A(B \to J/\psi K^*) = A_0 g_0 + A_\parallel g_\parallel + iA_\perp g_\perp, \tag{5}$$

$$\bar{A}(\bar{B} \to J/\psi \bar{K}^*) = \bar{A}_0 g_0 + \bar{A}_\parallel g_\parallel - i\bar{A}_\perp g_\perp, \tag{6}$$

with g_λ the coefficients of the helicity amplitudes in the linear polarization basis [10]. If one considers the case where K^* and $\bar{K}^*$ are detected through their decay to

$K_S\pi^0$ so that both B_d^0 and $\bar{B}_d^0$ decay to a common final state, the time-dependent decay rates can be written as

$$\Gamma[B_d(\bar{B}_d) \to J/\psi K^*] = e^{-\Gamma t}\sum_{\lambda\leq\sigma}(\Lambda_{\lambda\sigma} \pm \Sigma_{\lambda\sigma}\cos\Delta mt \mp \varrho_{\lambda\sigma}\sin\Delta mt)g_\lambda g_\sigma. \quad (7)$$

By performing an angular analysis and time-dependent study of the decays $\bar{B}_d \to J/\psi\bar{K}^*$ and $B_d \to J/\psi K^*$, one can measure the observables $\Lambda_{\lambda\sigma}$, $\Sigma_{\lambda\sigma}$, and $\varrho_{\lambda\sigma}$ [8]. These observables can be expressed in terms of the normalized helicity amplitudes A_0, $A_\|$, and $A_\perp$:

$$\Lambda_{\lambda\lambda} = \frac{|A_\lambda| + |\bar{A}_\lambda|^2}{2}, \quad \Sigma_{\lambda\lambda} = \frac{|A_\lambda|^2 - |\bar{A}_\lambda|^2}{2}, \quad \Lambda_{\perp i} = -\mathrm{Im}(A_\perp A_i^* - \bar{A}_\perp\bar{A}_i^*), \quad \Lambda_{\|0} = \mathrm{Re}(A_\| A_0^* + \bar{A}_\|\bar{A}_0^*),$$
$$\Sigma_{\perp i} = -\mathrm{Im}(A_\perp A_i^* + \bar{A}_\perp\bar{A}_i^*), \quad \Sigma_{\|0} = \mathrm{Re}(A_\| A_0^* - \bar{A}_\|\bar{A}_0^*), \quad \varrho_{\perp i} = -\mathrm{Re}\Big[\frac{q}{p}(A_\perp^*\bar{A}_i + A_i^*\bar{A}_\perp)\Big], \quad (8)$$
$$\varrho_{\perp\perp} = -\mathrm{Im}\Big(\frac{q}{p}A_\perp^*\bar{A}_\perp\Big), \quad \varrho_{\perp\perp} = \mathrm{Im}\Big[\frac{q}{p}(A_\|^*\bar{A}_0 + A_0^*\bar{A}_\|)\Big], \quad \varrho_{ii} = \mathrm{Im}\Big(\frac{q}{p}A_i^*\bar{A}_i\Big),$$

where $i = \{0, \|\}$, $q/p = \exp(-2i\phi_{\mathrm{mix}})$ with ϕ_{mix} the weak phase in $B_d^0 - \bar{B}_d^0$ mixing. From Eqs. (3) and (4) one can obtain the same observables in terms of a_λ, b_λ, ϕ, $\delta_\lambda \equiv \delta_\lambda^b - \delta_\lambda^a$, and $\Delta_i \equiv \delta_\perp^b - \delta_i^a$ [8]. In particular, $\Lambda_{\perp i}$ can be expressed as

$$\Lambda_{\perp i} = 2[a_\perp b_i\cos(\Delta_i - \delta_i) - a_i b_\perp\cos(\Delta_i + \delta_\perp)]\sin\phi. \quad (9)$$

The observable $\Lambda_{\perp i}$ is special, as made clear by Eq. (9), because it remains nonzero in the presence of NP effects ($\phi \neq 0$), even if the strong phase differences vanish. In contrast, direct CP asymmetries $\Sigma_{\lambda\lambda}$ are washed out if the strong phase differences vanish [8].

III. $B \to J/\psi K^*$

We can now proceed towards NP effects in $B \to J/\psi K^*$. We start by writing the decay amplitudes for $B \to J/\psi K^*$ using the factorization approximation, but keeping the color octet contribution [11],

$$\begin{aligned} A_\lambda(B\to J/\psi K^*) = {} & i\frac{G_F}{\sqrt{2}}V_{cb}V_{cs}^*f_\psi m_\psi\varepsilon_\psi^\mu(\lambda) \\ & \times\{a_2^{\mathrm{eff}}\langle K^*|\bar{s}\gamma_\mu(1-\gamma_5)b|B\rangle \\ & -\frac{\alpha_s}{2\pi}\frac{m_b}{q^2}\xi_8'\langle K^*|\bar{s}i\sigma_{\mu\nu}q^\nu[c_{12}\gamma_\mu(1+\gamma_5) \\ & + c_{12}'(1-\gamma_5)]b|B\rangle\}, \end{aligned} \quad (10)$$

where $\lambda = \lambda_{J/\psi} = \lambda_{K^*} = 0, \pm 1$ denote the helicities of the final state vector particles J/ψ and K^* in the B^0 rest frame [12]. The dominant contribution in Eq. (10) is given by the tree-level term proportional to

$$a_2^{\mathrm{eff}} = c_2^{\mathrm{eff}} + \zeta c_1^{\mathrm{eff}}, \quad (11)$$

while the color dipole moment terms, with the c_{12} operator coming dominantly from SM and the c_{12}' operator due exclusively to NP, give smaller corrections. In the expression for a_2^{eff} we have neglected the strong and electroweak penguin contributions. The quantity $\zeta = 1/N_c + \xi_8$ in Eq. (11) takes the value $1/N_c = 1/3$ in the naive factorization while deviations from $1/N_c$ due to nonfactorizable contributions to the hadronic matrix elements are measured by the parameters ξ_8 and ξ_8' [11]. The effective Wilson coefficients $c_1^{\mathrm{eff}}(m_b)$ and $c_2^{\mathrm{eff}}(m_b)$ for a $b \to s$ transition are defined in Ref. [13].

The way we proceed to determine the parameter ζ follows Ref. [14]. We fit the branching ratios for the decays $B \to J/\psi K_S(K^*)$ and $B \to \psi(2S)K^*$ to extract a_2^{eff}. We checked explicitly that the NP effect of Ref. [5] does not make significant impact on the decay rates. However, the extraction of a_2^{eff} depends on the specific model one uses for the hadronic form factors [14]. In this work we use the form factors at zero momentum transfer for the $B \to V$ transitions obtained in the light-cone sum rule (LCSR) analysis [15]. The form factor q^2 dependence is parametrized by

$$f(q^2) = \frac{f(0)}{1 - a(q^2/m_B^2) + b(q^2/m_B^2)^2}, \quad (12)$$

where the values of the parameters a and b are given in Ref [15]. Our extracted value for a_2^{eff} is 0.20 which, using Eq. (11) and the effective Wilson coefficients of Table 1 of Ref. [13], gives $\zeta = 0.48$, or the effective number of colors $1/\zeta = 2.1$. Knowing ζ one derives $\xi_8 = 0.15$ and for the value of ξ_8' we assume $\xi_8' = \xi_8$ [11].

IV. $B \to J/\psi K^*$ AND NP EFFECTS

As previously mentioned, in this work we focus on the decay $B_d^0(t) \to J/\psi K^*$ in the context of a supersymmetric model with maximal $\tilde{s}_R - \tilde{b}_R$ mixing. An approximate Abelian flavor symmetry [16] can be introduced to justify such a large mixing. In this model the right-right mass matrix $\tilde{M}_{RR}^{2(sb)}$ for the strange-beauty squark sector takes the form

$$\tilde{M}_{RR}^{2(sb)} = \begin{bmatrix} \tilde{m}_{22}^2 & \tilde{m}_{23}^2 e^{-i\sigma} \\ \tilde{m}_{23}^2 e^{i\sigma} & \tilde{m}_{33}^2 \end{bmatrix} \equiv R\begin{bmatrix} \tilde{m}_1^2 & 0 \\ 0 & \tilde{m}_2^2 \end{bmatrix}R^\dagger, \quad (13)$$

where $\tilde{m}_{ij} \sim \tilde{m}^2$, the squark mass scale, and

$$R = \begin{bmatrix} \cos\theta & \sin\theta \\ -\sin\theta e^{i\sigma} & \cos\theta e^{i\sigma} \end{bmatrix}. \tag{14}$$

The phase σ is the NP weak phase which will affect the observables one can extract from an angular analysis. Because of the almost democratic structure of $\tilde{M}^{2(sb)}_{RR}$, one of the two strange-beauty squarks, $\tilde{sb}_1$, can be rather light, even for $\tilde{m} \sim O(\text{TeV})$. A light strange-beauty squark, together with a light gluino, can make S_{ϕ,K_s} negative for $\sigma \lesssim \pi/2$ [5], and one still survives [5,17] the usually stringent $b \to s\gamma$ constraint.

The main new contribution to $B^0_d(t) \to J/\psi K^*$ is given by the color dipole moment amplitude through gluino and $\tilde{sb}$ squark exchange in the loop. The analytic expressions for the Wilson coefficient c'_{12} for the color dipole moment operator $g_s/(8\pi^2)m_b\bar{s}_\alpha\sigma^{\mu\nu}(1 \pm g_5)\lambda^A_{\alpha\beta}/2b_\beta G^A_{\mu\nu}$ can be found in Ref. [18]. The coefficient c_{12} remains basically the same as $c^{\text{SM}}_{12} = -0.15$. In Fig. 1 we plot the real and imaginary parts of the Wilson coefficient c'_{12} for three different values of gluino mass, $m_{\tilde{g}} = 300, 500, 800$ GeV [19]. The eigenvalues of Eq. (13) are taken (with some level of tuning) as $\tilde{m}^2_1 = (200\text{ GeV})^2$ and $\tilde{m}^2_2 = 2\tilde{m}^2 - \tilde{m}^2_1$ with $\tilde{m}^2 = (2\text{ TeV})^2$.

Using the parametrization for the matrix elements $\langle K^* \mid \bar{s}\gamma_\mu(1-\gamma_5)b \mid B\rangle$ and $\langle K^* \mid \bar{s}i\sigma_{\mu\nu}\gamma_\mu(1\pm\gamma_5)b \mid B\rangle$ given in Ref. [11] the amplitudes $A_\lambda(B \to J/\psi K^*)$ can be written as

$$\begin{aligned} A_\lambda(B \to J/\psi K^*) = {} & i\frac{G_F}{\sqrt{2}}V_{cb}V^*_{cs}f_\psi m_\psi \epsilon^*_{K^*\mu}(\lambda)\epsilon^*_{\psi\nu}(\lambda)\Bigg(a^{\text{eff}}_2\Bigg[\frac{V(m^2_\psi)}{m_B + m_{K^*}}\varepsilon^{\mu\nu\alpha\beta}p_{K^*\alpha}p_{B\beta} - \frac{i}{2}(m_B + m_{K^*})A_1(m^2_\psi)g^{\mu\nu} \\ & + i\frac{A_2(m^2_\psi)}{m_B + m_{K^*}}p^\mu_B p^\nu_B\Bigg] - \frac{\alpha_s}{2\pi}\frac{m_B}{m^2_\psi}\xi_8(c_{12} + c'_{12})g_+(m^2_\psi)\varepsilon^{\mu\nu\alpha\beta}p_{K^*\alpha}p_{B\beta} + i\frac{\alpha_s}{2\pi}\frac{m_B}{m^2_\psi}\xi_8(c_{12} - c'_{12}) \\ & \times\left\{\frac{1}{2}[g_+(m^2_\psi)(m^2_B - m^2_{K^*}) + g_+(m^2_\psi)m^2_\psi]g^{\mu\nu} - [g_+(m^2_\psi) - h(m^2_\psi)m^2_\psi]p_{B\mu}p_{B\nu}\right\}\Bigg), \end{aligned} \tag{15}$$

The general covariant form for $A_{0,\pm1}$ is given by

$$\begin{aligned} A_{0,\pm1} = {} & \epsilon^*_{\psi\mu}(0,\pm1)\epsilon^*_{K^*\nu}(0,\pm1)\Bigg[ag^{\mu\nu} + \frac{b}{m_\psi m_{K^*}}p^\mu_B p^\nu_B \\ & + \frac{ic}{m_\psi m_{K^*}}\epsilon^{\mu\nu\alpha\beta}p_{K^*\alpha}p_{B\beta}\Bigg], \end{aligned} \tag{16}$$

where a, b, and c are three invariant amplitudes. The corresponding amplitudes $\bar{A}_{0,\pm1}$ are obtained by taking the conjugate of the invariant amplitudes a, b, c and switching the sign of the term $\epsilon^{\mu\nu\alpha\beta}p_{K^*\alpha}p_{B\beta}$ in Eq. (16). By comparing Eq. (15) with Eq. (16) one can extract the three invariant amplitudes,

$$\begin{aligned} a = {} & \frac{1}{2}(m_B + m_{K^*})A_1(m^2_\psi)a^{\text{eff}}_2 - \frac{\alpha_s}{2\pi}\frac{m_B}{m^2_\psi}\xi_8(c_{12} - c'_{12}) \\ & \times\frac{1}{2}[g_+(m^2_\psi)(m^2_B - m^2_{K^*}) + g_+(m^2_\psi)m^2_\psi], \end{aligned} \tag{17}$$

$$\begin{aligned} b = {} & -m_\psi m_{K^*}\Bigg\{\frac{A_2(m^2_\psi)}{m_B + m_{K^*}}a^{\text{eff}}_2 + \frac{\alpha_s}{2\pi}\frac{m_B}{m^2_\psi}\xi_8(c_{12} - c'_{12}) \\ & \times[g_+(m^2_\psi) - h(m^2_\psi)m^2_\psi]\Bigg\}, \end{aligned} \tag{18}$$

$$\begin{aligned} c = {} & m_\psi m_{K^*}\Bigg[\frac{V(m^2_\psi)}{m_B + m_{K^*}}a^{\text{eff}}_2 - \frac{\alpha_s}{2\pi}\frac{m_B}{m^2_\psi}\xi_8(c_{12} \\ & + c'_{12})g_+(m^2_\psi)\Bigg]. \end{aligned} \tag{19}$$

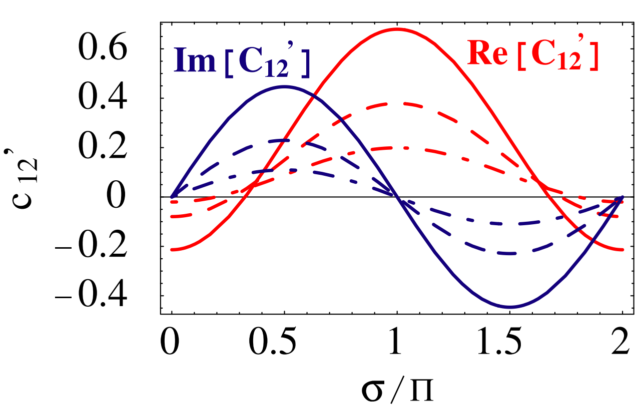

FIG. 1 (color online). c'_{12} as a function of σ for $\tilde{m}_1 = 200$ GeV and $\tilde{m} = 2$ TeV. $\text{Re}(c'_{12})$ and $\text{Im}(c'_{12})$ are plotted with solid, dashed, and dot-dashed lines for $m_{\tilde{g}} = 300, 500, 800$ GeV.

In the expressions of Eqs. (17)–(20) we have omitted the common factor $\sqrt{2}G_F V_{cb}V^*_{cs}f_\psi m_\psi$. Note that each invariant amplitude contains a SM contribution which is dominated by the tree-level term proportional to a^{eff}_2 plus the color dipole moment term proportional to c_{12}, and a NP contribution proportional to c'_{12}. One can consequently write the three invariant amplitudes as the sum of a SM and a NP contribution: $a, b, c \equiv (a, b, c)^{\text{SM}} + (a, b, c)^{\text{NP}}$.

We now rewrite the helicity amplitudes $A_{0,\pm1}$ in terms of the invariant amplitudes and the kinematic factor $x \equiv p_\psi \cdot$

$p_{K^*}/(m_\psi m_{K^*})$ [20]. Writing separately for the SM and NP contributions, we obtain

$$A_{\pm1}^{\rm SM,NP} = a^{\rm SM,NP} \pm c^{\rm SM,NP}\sqrt{x^2-1},$$
$$A_0^{\rm SM,NP} = -a^{\rm SM,NP}x - b^{\rm SM,NP}(x^2-1), \qquad (20)$$

with $A_{0,\pm1} = A_{0,\pm1}^{SM} + A_{0,\pm1}^{NP}$. To evaluate the observables of Eq. (8) one can transform $A_{0,\pm1}^{\rm SM,NP}$ in the corresponding linear polarization amplitudes $A_{0,\parallel,\perp}^{\rm SM,NP}$ using the relations: $A_{\parallel,\perp} = (A_{+1} \pm A_{-1})/\sqrt{2}$, A_0 being the same in both bases.

The invariant amplitudes in the linear polarization basis $a_{0,\parallel,\perp}$ and $b_{0,\parallel,\perp}$ can subsequently be expressed in terms of $(a,b,c)^{\rm (SM,NP)}$,

$$a_0 = -a^{\rm SM}x - b^{\rm SM}(x^2-1),$$
$$b_0 e^{i\phi} = -a^{\rm NP}x - b^{\rm NP}(x^2-1), \qquad a_\parallel = \sqrt{2}a^{\rm SM},$$
$$b_\parallel e^{i\phi} = \sqrt{2}a^{\rm NP}, \qquad a_\perp = c^{\rm SM}\sqrt{2(x^2-1)}, \qquad (21)$$
$$b_\perp e^{i\phi} = c^{\rm NP}\sqrt{2(x^2-1)},$$

where $\phi = \arg(c'_{12}) \pm \pi$.

In Fig. 2 we plot the observables $\Lambda_{\perp i}$ versus the NP weak phase $\sigma = \arg(c'_{12})$ for three different values of the gluino mass, $m_{\tilde g} = 300, 500, 800$ GeV, with squark masses $\tilde m_1^2 = (200\text{ GeV})^2$ and $\tilde m_2^2 = 2\tilde m^2 - \tilde m_1^2$ with $\tilde m = 2$ TeV. We note that the effects of NP can be at most a few percent, and tend to disappear as the gluino becomes heavier [21]. This can be understood from Eq. (9) together with the expressions for $\Lambda_{\lambda\lambda} = a_\lambda^2 + b_\lambda^2 + 2a_\lambda b_\lambda \cos\delta_\lambda \cos\phi$. In fact, the main contributions to $\Lambda_{\lambda\lambda}$ are proportional to a_λ^2 (tree-level dominated) hence are of $O(1)$, while for $\Lambda_{\perp i}$ one has $O(b_\lambda/a_\lambda) \sim O(0.01)$ at most, being suppressed by the ratio $P^{\rm NP}/T^{\rm SM}$ with P and T indicating, respectively, the penguin and tree terms.

We see that, for the particular model considered in this work, to observe NP effects by performing an angular analysis for the decay $B_d^0(t) \to J/\psi K^*$, one needs to be able to extract $\Lambda_{\perp i}$ with a precision of at least a few percent. On the other hand, as stressed in Ref. [8], no tagging or time-dependent measurements are needed to measure $\Lambda_{\perp i}$ since it appears with the same sign in both rates for $B_d^0(t)$ and $\bar B_d^0(t)$ (see Eq. (7)).

Following from the above considerations, it is evident that the decay $B_d^0(t) \to \phi K^*$ becomes really interesting. This decay, contrary to $B_d^0(t) \to J/\psi K^*$, is not tree-level dominated. Rather, it is of pure penguin type, and in the model considered a_λ and b_λ are of the same order. This implies now that the observables $\Lambda_{\perp i}$ not only differ from zero if there are NP effects, but they are expected to be of $O(1)$, with the ratio $P^{\rm NP}/P^{\rm SM} \sim O(1)$. Obviously for the decay $B_d^0(t) \to \phi K^*$ the hadronic uncertainties play a more important role than for the decay $B_d^0(t) \to J/\psi K^*$, plaguing the theoretical prediction for $\Lambda_{\perp i}$. Furthermore, the strength of the transverse components are not yet understood.

For $B_d^0(t) \to J/\psi K^*$ decay, full angular analysis by the CLEO Collaboration [22] shows that the P-wave component is small, $|P|^2 = |A_\perp|^2 = 0.16 \pm 0.08 \pm 0.04$, while the longitudinal component is around 50%, $\Gamma_L/\Gamma = |A_0|^2 = 0.52 \pm 0.07 \pm 0.04$. Recent measurements for the longitudinal and transverse amplitudes have been also reported by both *BABAR* [23] and Belle [24] collaborations with the respective values $|A_\perp|^2 = 0.16 \pm 0.03 \pm 0.01$, $|A_0|^2 = 0.60 \pm 0.03 \pm 0.02$ and $|A_\perp|^2 = 0.19 \pm 0.02 \pm 0.03$, $|A_0|^2 = 0.62 \pm 0.02 \pm 0.03$. In Figs. 3 and 4 we plot respectively $|P|^2$ and Γ_L/Γ as a function of the NP weak phase σ. It seems clear that both statistical and systematic errors need to be reduced by an order of magnitude to discriminate a nonzero value for $\Lambda_{\perp i}$ as predicted by the model considered. We expect that the error on the extracted value of $\Lambda_{\perp i}$ will be of the same order as the one on $|P|^2$ or Γ_L/Γ.

We conclude this section by presenting the results for $\sin 2\phi_{1\lambda}^{\rm meas} = {\rm Im}[(q/p)(\bar A_\lambda/A_\lambda)]$, where $\lambda = 0, \parallel, \perp$ [25].

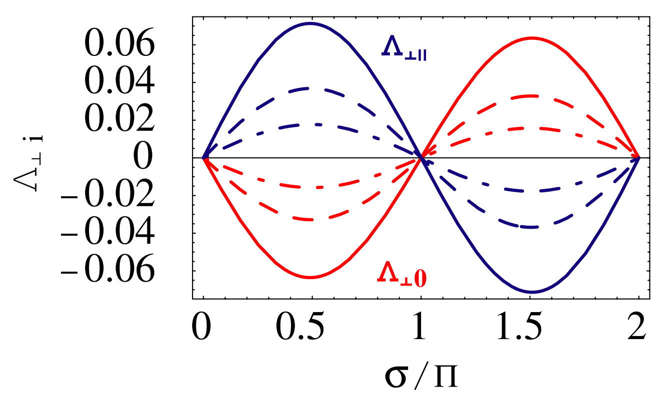

FIG. 2 (color online). $\Lambda_{\perp i}$ as a function of σ for $\tilde m_1 = 200$ GeV and $\tilde m = 2$ TeV. $\Lambda_{\perp i}$ are plotted in solid, dashed, and dot-dashed lines for $m_{\tilde g} = 300, 500, 800$ GeV.

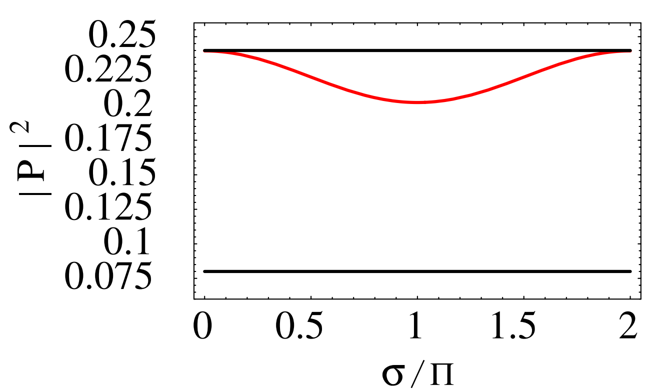

FIG. 3 (color online). $|P|^2$ as a function of σ for $\tilde m_1 = 200$ GeV, $\tilde m = 2$ TeV, and $m_{\tilde g} = 300$ GeV compared with the experiment.

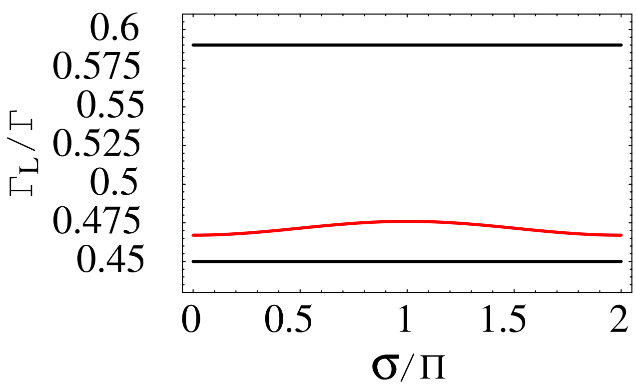

FIG. 4 (color online). Γ_L/Γ as a function of σ for $\tilde{m}_1 = 200$ GeV, $\tilde{m} = 2$ TeV, and $m_{\tilde{g}} = 300$ GeV compared with the experiment.

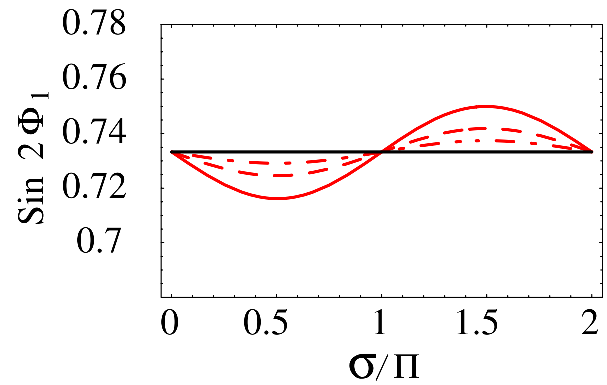

FIG. 6 (color online). $\sin 2\phi_1(B \to J/\psi K_S)$ as a function of σ for $\tilde{m}_1 = 200$ GeV and $\tilde{m} = 2$ TeV. $\sin 2\phi_1(B \to J/\psi K_S)$ is plotted in solid, dashed, and dot-dashed lines for $m_{\tilde{g}} = 300$, 500, 800 GeV. The solid horizontal line corresponds to $\sin 2\phi_1^{\text{mix}} = 0.733$.

If NP effects are present one will not be able to extract ϕ_1^{mix} [26] (the phase of $B_d^0 - \bar{B}_d^0$ mixing which in general can be affected by NP) and the measured value of ϕ_1, which will depend on the helicity of the final state, will differ from the real value of ϕ_1^{mix} [8]. In Fig. 5 we plot $\sin 2\phi_{1\lambda}^{\text{meas}}$ as a function of the NP weak phase σ. As for the quantities $\Lambda_{\perp i}$, the effects of NP on $\sin 2\phi_{1\lambda}^{\text{meas}}$ can be at most a few percent. Deviations from $\sin 2\phi_1^{\text{mix}}$ reach their largest value at $\sigma = \pi/2$ and are bigger for the transverse components, $\{\lambda = \|, \perp\}$. On the other hand deviations on $\sin 2\phi_{10}^{\text{meas}}$, even if smaller, could be easier to detect because of the higher number of longitudinally polarized final states. On top of that, by comparing $\sin 2\phi_1(B \to J/\psi K_S)$ [11] in Fig. 6 with $\sin 2\phi_{10}^{\text{meas}}$, one can observe that deviations from $\sin 2\phi_1^{\text{mix}}$ have opposite signs [27]. This divergent behavior could in principle be easier to observe than the single deviations.

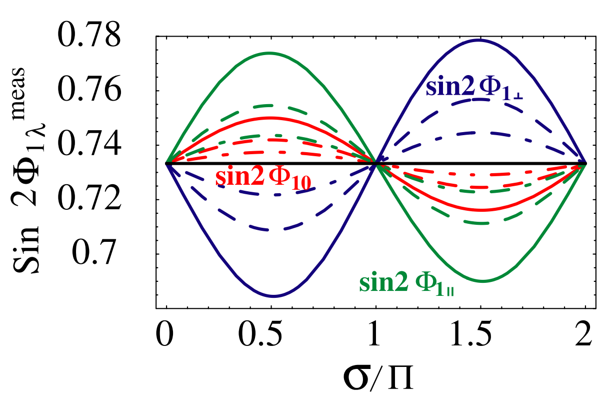

FIG. 5 (color online). $\sin 2\phi_{1\lambda}^{\text{meas}}$ as a function of σ for $\tilde{m}_1 = 200$ GeV and $\tilde{m} = 2$ TeV. $\sin 2\phi_{1\lambda}^{\text{meas}}$ are plotted in solid, dashed, and dot-dashed lines for $m_{\tilde{g}} = 300$, 500, 800 GeV. The solid horizontal line corresponds to $\sin 2\phi_1^{\text{mix}} = 0.733$.

V. CONCLUSION

A supersymmetric extension of the SM with a light right-handed "strange-beauty" squark, a light gluino, and a new *CP* phase seems to contain all the necessary ingredients to explain the recent *CP* anomaly in $B_d \to \phi K_S$. In the same framework we have calculated possible NP effects to observables that can be extracted by the time-dependent angular analysis of $B_d \to J/\psi K^*$. An important role is played by the quantities $\Lambda_{\perp i}$ with $i = \{0, \|\}$, which can be nonzero in the presence of NP even for very small strong phase differences. Our results show that deviations from zero can be at most of the order of a few percent, since it is suppressed by the ratios of the NP penguin amplitude to the SM tree amplitude. This obviously suggests that for decays which are pure penguins, like $B_d \to \phi K^*$, deviations from zero for the observables $\Lambda_{\perp i}$ are expected to be of order 1.

The quantities $\sin 2\phi_{1\lambda}^{\text{meas}}$ can also differ from the real value $\sin 2\phi_1^{\text{mix}}$ because of NP effects. In particular, $\sin 2\phi_{10}^{\text{meas}}$ and $\sin 2\phi_1(B \to J/\psi K_S)$ have opposite deviations from $\sin 2\phi_1^{\text{mix}}$, and by comparing the two behaviors, NP effects could be easier to observe than by looking at single deviations. As for $\Lambda_{\perp i}$, we found that deviations are of the order of a few percent.

In conclusion, new physics effects to $B_d \to J/\psi K^*$ from the model considered in this work are found to be too small to be observed at the current B factories. But because of the small NP effects, $B_d^0(t) \to J/\psi K^*$ remains a good mode for measuring $\sin 2\phi_1^{\text{mix}}$.

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[1] *BABAR* Collaboration, B. Aubert *et al.*, hep-ex/0408076.
[2] Belle Collaboration, K. Abe *et al.*, hep-ex/0409049.
[3] Moriond 2004, http://www.slac.stanford.edu/xorg/hfag/index.html.
[4] M. Ciuchini, E. Franco, G. Martinelli, A. Masiero, and L. Silvestrini, Phys. Rev. Lett. **79**, 978 (1997); E. Lunghi and D. Wyler, Phys. Lett. B **521**, 320 (2001).
[5] C. K. Chua, W.-S. Hou, and M. Nagashima, Phys. Rev. Lett. **92**, 201803 (2004).
[6] X.-G. He and W.-S. Hou, Phys. Rev. D **58**, 117502 (1998).
[7] R. Fleischer, Phys. Rev. D **60**, 073008 (1999).
[8] D. London, N. Sinha, and R. Sinha, Phys. Rev. Lett. **85**, 1807 (2000); A. Datta and D. London, Int. J. Mod. Phys. A **19**, 2505 (2004); D. London, N. Sinha, and R. Sinha, Europhys. Lett. **67**, 579 (2004); hep-ph/0307308; hep-ph/0207007; Phys. Rev. D **69**, 114013 (2004).
[9] J. L. Rosner, Phys. Rev. D **42**, 3732 (1959).
[10] N. Sinha and R. Sinha, Phys. Rev. Lett. **80**, 3706 (1998).
[11] X.-G. He and W.-S. Hou, Phys. Lett. B **445**, 344 (1999).
[12] G. Kramer and W. F. Palmer, Phys. Rev. D **45**, 193 (1992).
[13] Y.-H. Chen, H.-Y. Cheng, B. Tseng, and K.-C. Yang, Phys. Rev. D **60**, 094014 (1999).
[14] M. Neubert and B. Stech, Adv. Ser. Dir. High Energy Phys. **15**, 294 (1998).
[15] P. Ball and V. M. Braun, Phys. Rev. D **58**, 094016 (1998).
[16] Y. Nir and Sieberg, Phys. Lett. B **309**, 337 (1993); M. Leurer, Y. Nir, and N. Seiberg, Nucl. Phys. **B420**, 468 (1994).
[17] A. Arhrib, C.-K. Chua, and W.-S. Hou, Phys. Rev. D **65**, 017701 (2002).
[18] A. Arhrib, C.-K. Chua, and W.-S. Hou, Eur. Phys. J. C **21**, 567 (2001).
[19] A light gluino of 300 GeV is disfavored by the $b \to s\gamma$ constraint [5,17].
[20] A. S. Dighe, I. Dunietz, and J. L. Roner, Phys. Lett. B **369**, 144 (1996).
[21] Conservative variations of the squark scale $\tilde{m}$ and the squark mass $\tilde{m}_1$ do not produce valuable changes in the observables $\Lambda_{\perp i}$ from the case presented in Fig. 2.
[22] CLEO Collaboration, C. P. Jessop *et al.*, Phys. Rev. Lett. **79**, 4533 (1997).
[23] *BABAR* Collaboration, B. Aubert *et al.*,Phys. Rev. Lett.**87**, 241801 (2001).
[24] Belle Collaboration, K. Abe *et al.*, Phys. Lett. B **538**, 11 (2002).
[25] The available measurements of $\sin 2\phi_1$ from *BABAR* Collaboration with the value 0.22 ± 0.52 and Belle Collaboration with the value 0.10 ± 0.45 [3] are obtained without separating final states of different *CP*.
[26] The number of measurements for the decays $B \to V_1 V_2$ and $\bar{B} \to \bar{V}_1 \bar{V}_2$ is fewer than the number of theoretical parameters, making it impossible to predict $\phi_1^{\rm mix}$ purely in terms of observables.
[27] The final state for $B \to J/\psi K_S$ is *CP* odd on the contrary to the longitudinal component for $B \to J/\psi K^*$ with $K^* \to K_S \pi^0$ which is *CP* even.