Possible resolution of the B-meson decay polarization anomaly in R-parity violating supersymmetry

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(Received 2 February 2007; published 27 April 2007)

We examine the possible resolution of the recently observed polarization anomaly in $B^0(\bar{B}^0) \to \phi K^{*0}(\bar{K}^{*0})$-decay within R-parity violating ($\not{R}_p$) SUSY. We show that a combination of the superpotential trilinear $\not{R}_p$-interactions, with the couplings λ', and the soft SUSY breaking bilinear $\not{R}_p$ sneutrino-Higgs mixing, proportional to $\tilde{\mu}^2$, can potentially generate the effective operators with the chirality structure necessary to account for this anomaly. However, we demonstrate that the existing experimental data on $B_s \to \mu^+\mu^-$-decay lead to stringent upper limits on the Wilson coefficients of these operators, which are about 2 orders of magnitude below the values required for the resolution of the B-decay polarization anomaly, and, therefore, it can hardly be explained within the $\not{R}_p$ SUSY framework. As a byproduct result of our analysis we derive new limits on the products of the soft bilinear and the superpotential trilinear $\not{R}_p$-parameters of the form $\tilde{\mu}^2\lambda'$.

DOI: 10.1103/PhysRevD.75.074029 PACS numbers: 12.39.Fe, 11.30.Er, 13.40.Em, 14.20.Dh

I. INTRODUCTION

Now it is widely recognized that B-mesons offer powerful means for testing the standard model (SM) and probe physics beyond its framework. Recently, remarkable progress has been achieved in experimental and theoretical studies of B-physics. One of the most important experimental results of the last years in this field was, certainly, the discovery of CP violation in the B-system. The running B-experiments [1] at *BABAR*, BELLE, CDF, D0 and CLEO have also collected a large statistics on various decay modes of B-mesons some of which seem to be quite challenging for the SM.

The *BABAR* [2] and BELLE [3] Collaborations reported experimental data on B-meson decay to a pair of light vector mesons: $B \to VV$ where $V = \rho, \phi, K^*$. An intriguingly large transverse polarization fraction comparable to the longitudinal one has been observed in the $B^0(\bar{B}^0) \to \phi K^{*0}(\bar{K}^{*0})$-decay channel. This result has been recently confirmed by the CDF collaboration [4] as well. This polarization anomaly is hard to be explained within the SM and may indicate some new physics. As is known, the SM predicts for the helicity amplitudes of $B^0 \to \phi K^{*0}$ the following ratios [5,6]: $H_{00}:H_{--}:H_{++} \sim \mathcal{O}(1):\mathcal{O}(1/m_b):\mathcal{O}(1/m_b^2)$, where H_{00} corresponds to the final vector mesons in the longitudinal polarization state, while H_{++}, H_{--} in the transverse positive and negative one. This SM result is in an obvious disagreement with the *BABAR* [2], BELLE [3] and CDF [4] observations, demonstrating that $|H_{++} \pm H_{--}|^2 \approx |H_{00}|^2$.

In the literature various efforts have been undertaken to account for the polarization anomaly from the view point of the SM [6,7] and in various scenarios of new physics beyond the SM [8,9]. In Ref. [9] a model independent analysis of the B-decay polarization anomaly has been carried out on the basis of the effective Lagrangian approach. Two sets of effective $\Delta B = 1$ operators necessary for the resolution of this anomaly have been identified. In addition from the experimental data [2,3] the corresponding values of their Wilson coefficients have been determined. These two sets of operators have the following chirality structure: (i) $(1-\gamma_5)\otimes(1-\gamma_5)$, $\sigma(1-\gamma_5)\otimes\sigma(1-\gamma_5)$ and (ii) $(1+\gamma_5)\otimes(1+\gamma_5)$, $\sigma(1+\gamma_5)\otimes\sigma(1+\gamma_5)$.

In the present paper we use this model independent result in order to examine the ability of R-parity violating SUSY ($\not{R}_p$ SUSY) to resolve the above discussed polarization anomaly in $B^0(\bar{B}^0) \to \phi K^{*0}(\bar{K}^{*0})$-decay. In Sec. II we specify the effective $\Delta B = 1$ operators satisfying the polarization Anomaly Resolution Criteria (pARC). In Sec. III we determine these operators in the context of $\not{R}_p$ SUSY and derive their Wilson coefficients. In Sec. IV we study experimental limits on these Wilson coefficients from the existing $B_s \to \mu^+\mu^-$ data and discuss the compatibility of the pARC with these limits.

II. CRITERIA FOR RESOLUTION OF THE POLARIZATION ANOMALY

The effective Hamiltonian $\mathcal{H}$ describing $\bar{B}^0 \to \phi\bar{K}^{*0}$ with $\Delta B = 1$ can be written in the form

$$\mathcal{H}_{\Delta B=1} = \mathcal{H}^{\rm SM}_{\Delta B=1} + \mathcal{H}^{\rm NP}_{\Delta B=1} = \frac{G_F}{\sqrt{2}}\sum_{i=1}^{14} c_i(\mu)\cdot O_i(\mu) + \frac{G_F}{\sqrt{2}}\sum_{i=15}^{44} c_i(\mu)\cdot O_i(\mu) + \text{H.c.}, \quad (1)$$

where $c_i(\mu)$ are the Wilson coefficients evaluated at the

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renormalization scale $\mu \sim m_b$. The first 14 terms correspond to the penguin-dominated SM contributions $\mathcal{H}^{\rm SM}_{\Delta B=1}$ listed in Ref. [10], the last 30 terms $\mathcal{H}^{\rm NP}_{\Delta B=1}$ appear in the presence of new physics (NP). In Ref. [9] it was shown that out of the 30 NP-operators only the following operator set

$$O_{15} = \bar{s}_\alpha P_R b_\alpha \cdot \bar{s}_\beta P_R s_\beta, \qquad O_{16} = \bar{s}_\alpha P_R b_\beta \cdot \bar{s}_\beta P_R s_\alpha, \tag{2}$$

$$O_{17} = \bar{s}_\alpha P_L b_\alpha \cdot \bar{s}_\beta P_L s_\beta, \qquad O_{18} = \bar{s}_\alpha P_L b_\beta \cdot \bar{s}_\beta P_L s_\alpha, \tag{3}$$

$$\begin{aligned} O_{23} &= \bar{s}_\alpha \sigma^{\mu\nu} P_R b_\alpha \cdot \bar{s}_\beta \sigma_{\mu\nu} P_R s_\beta, \\ O_{24} &= \bar{s}_\alpha \sigma^{\mu\nu} P_R b_\beta \cdot \bar{s}_\beta \sigma_{\mu\nu} P_R s_\alpha, \end{aligned} \tag{4}$$

$$\begin{aligned} O_{25} &= \bar{s}_\alpha \sigma^{\mu\nu} P_L b_\alpha \cdot \bar{s}_\beta \sigma_{\mu\nu} P_L s_\beta, \\ O_{26} &= \bar{s}_\alpha \sigma^{\mu\nu} P_L b_\beta \cdot \bar{s}_\beta \sigma_{\mu\nu} P_L s_\alpha, \end{aligned} \tag{5}$$

satisfies the polarization Anomaly Resolution Criteria (pARC) [9], allowing one to possibly solve the polarization anomaly in $B^0(\bar{B}^0) \to \phi K^{*0}(\bar{K}^{*0})$-decay. Here, α and β are the color indices. In what follows we denote the set of operators in Eqs. (2)–(5) as pARC operators. In Ref. [9] it was noted that the (pseudo-)scalar operators O_{15-18} can be expressed in the basis of (pseudo-)tensor operators O_{23-26} by Fierz transformation

$$O_{15} = \tfrac{1}{12} O_{23} - \tfrac{1}{6} O_{24}, \qquad O_{16} = \tfrac{1}{12} O_{24} - \tfrac{1}{6} O_{23}, \tag{6}$$

$$O_{17} = \tfrac{1}{12} O_{25} - \tfrac{1}{6} O_{26}, \qquad O_{18} = \tfrac{1}{12} O_{26} - \tfrac{1}{6} O_{25}. \tag{7}$$

The contributions of the operators O_{15-26} to the helicity amplitudes of $\bar{B}^0 \to \phi \bar{K}^{*0}$-decay can be calculated within the QCD factorization (QCDF) approach in terms of the corresponding Wilson coefficients and hadronic form factors. In this approach the helicity amplitudes take the form [6,9]

$$\bar{H}_{00} = -4 i f^T_\phi m_B^2 (\tilde{a}_{23} - \tilde{a}_{25}) [h_2 T_2(m_\phi^2) - h_3 T_3(m_\phi^2)], \tag{8}$$

$$\begin{aligned} \bar{H}_{\pm\pm} = -4 i f^T_\phi m_B^2 \{ & \tilde{a}_{23} [\pm f_1 T_1(m_\phi^2) - f_2 T_2(m_\phi^2)] \\ & + \tilde{a}_{25} [\pm f_1 T_1(m_\phi^2) + f_2 T_2(m_\phi^2)] \}. \end{aligned} \tag{9}$$

Here the ϕ-meson tensor decay constant f^T_ϕ and the form factors of the $\bar{B} - \bar{K}^*$ transition are defined as

$$\langle \phi(q, \epsilon_1) | \bar{s} \sigma^{\mu\nu} s | 0 \rangle = -i f^T_\phi (\epsilon_1^{\mu *} q^\nu - \epsilon_1^{\nu *} q^\mu), \tag{10}$$

$$\begin{aligned} &\langle \bar{K}^*(p', \epsilon_2) | \bar{s} \sigma_{\mu\nu} q^\nu (1 + \gamma_5) s | \bar{B}(p) \rangle \\ &\quad = 2 i \epsilon_{\mu\nu\rho\sigma} \epsilon_2^{\nu *} p^\rho p'^\sigma T_1(q^2) + T_2(q^2) [\epsilon^*_{2\mu} (m_B^2 - m_{K^*}^2) \\ &\qquad - (\epsilon_2^* \cdot p)(p + p')_\mu] + T_3(q^2)(\epsilon_2^* \cdot p) \\ &\qquad \times \left[q_\mu - \frac{q^2 (p + p')_\mu}{m_B^2 - m_{K^*}^2} \right], \end{aligned} \tag{11}$$

with $q = p - p'$ and $m_B = 5.279$ GeV, $m_{K^*} = 0.892$ GeV, $m_\phi = 1.019$ GeV being the masses of the B, K^* and ϕ mesons, respectively. The kinematical factors in Eqs. (8) are

$$f_1 = \frac{2 p_c}{m_B}, \qquad f_1 = \frac{m_B^2 - m_{K^*}^2}{m_B^2}, \tag{12}$$

$$h_2 = \frac{1}{2 m_{K^*} m_\phi} \left[\frac{(m_B^2 - m_\phi^2 - m_{K^*}^2)(m_B^2 - m_{K^*}^2)}{m_B^2} - 4 p_c^2 \right], \tag{13}$$

$$h_3 = \frac{1}{2 m_{K^*} m_\phi} \frac{4 p_c^2 m_\phi^2}{m_B^2 - m_{K^*}^2}, \tag{14}$$

where p_c is the momentum of the ϕ or K^* meson in the rest frame of the decaying $\bar{B}^0$ meson. The effective coefficients in Eq. (8) are expressed in terms of the Wilson coefficients as [9]

$$\begin{aligned} \tilde{a}_{23} = {} & (1 + \tfrac{1}{2N_c})(c_{23} + \tfrac{1}{12} c_{15} - \tfrac{1}{6} c_{16}) + (\tfrac{1}{N_c} + \tfrac{1}{2})(c_{24} \\ & + \tfrac{1}{12} c_{16} - \tfrac{1}{6} c_{15}) + \text{nonfact.}, \end{aligned} \tag{15}$$

$$\begin{aligned} \tilde{a}_{25} = {} & (1 + \tfrac{1}{2N_c})(c_{25} + \tfrac{1}{12} c_{17} - \tfrac{1}{6} c_{18}) + (\tfrac{1}{N_c} + \tfrac{1}{2})(c_{25} \\ & + \tfrac{1}{12} c_{18} - \tfrac{1}{6} c_{17}) + \text{nonfact.} \end{aligned} \tag{16}$$

The last terms in (15) and (16) indicate corrections due to deviations from the QCDF.

On the basis of the above equations in Ref. [9] there has been made an analysis of the experimental data obtained by *BABAR* [2] and BELLE [3] on the angular distribution in $B^0(\bar{B}^0) \to \phi K^{*0}(\bar{K}^{*0})$-decay. It was shown that there are two theoretical scenarios which can separately account for the polarization anomaly of these data.

Scenario (i) $\tilde{a}_{23} = 0$ and

$$\begin{aligned} |\tilde{a}_{25}| = 2.10^{+0.19}_{-0.12} \times 10^{-4}, \qquad & \delta_{25} = 1.15 \pm 0.09, \\ \phi_{25} = -0.12 \pm 0.09. & \end{aligned} \tag{17}$$

Scenario (ii) $\tilde{a}_{25} = 0$ and

$$\begin{aligned} |\tilde{a}_{23}| = 1.70^{+0.11}_{-0.07} \times 10^{-4}, \qquad & \delta_{23} = 2.36 \pm 0.10, \\ \phi_{23} = 0.14 \pm 0.09. & \end{aligned} \tag{18}$$

Here the following notations were used $\tilde{a}_{ij} = |\tilde{a}_{ij}| e^{i\delta_{ij}} e^{i\phi_{ij}}$, identifying ϕ_{ij} and δ_{ij} with the weak (coming from the terms in Eq. (1)) and strong phases, respectively. The values of Eqs. (17) and (18) correspond to the best fit values for the combined data of *BABAR* [2] and BELLE [3]. In what follows we use these results as a criterion to assess if a particular model is able to resolve the polarization anomaly in question or not.

III. PARC OPERATORS IN $\not{R}_p$ SUSY

In relation to the polarization anomaly in $\bar{B}^0 \to \phi \bar{K}^{*0}$ we are studying the $\Delta B = 1$ transitions on the quark level. Here we derive the effective Lagrangian describing these transitions within the minimal $\not{R}_p$ SUSY model ($\not{R}_p$ MSSM) and show that among the resulting set of operators there emerge the pARC operators O_{15} and O_{17}. In the generic case of $\not{R}_p$ MSSM *R*-parity is violated by the following terms in the superpotential

$$W_{\not{R}_p} = \mu_j L_j H_2 + \tfrac{1}{2}\lambda_{ijk} L_i L_j E^c_k + \bar{\lambda}'_{ijk} L_i Q_j D^c_k + \tfrac{1}{2}\bar{\lambda}''_{ijk} U^c_i D^c_j D^c_k, \tag{19}$$

and in the soft SUSY breaking part of the scalar potential

$$V^{\rm soft}_{\not{R}_p} = \Lambda_{ijk}\tilde{L}_i\tilde{L}_j\tilde{E}^c_k + \Lambda'_{ijk}\tilde{L}_i\tilde{Q}_j\tilde{D}^c_k + \Lambda''_{ijk}\tilde{U}^c_i\tilde{D}^c_j\tilde{D}^c_k + \tilde{\mu}^2_{2j}\tilde{L}_j H_2 + \tilde{\mu}^2_{1j}\tilde{L}_j H_1^\dagger + \text{H.c.} \tag{20}$$

In Eq. (19) L, Q stand for the lepton and quark doublet left-handed superfields, while E^c, U^c, D^c for the lepton and *up*, *down* quark singlet superfields; H_2 is the Higgs doublet superfields with a weak hypercharge $Y = +1$, respectively. In Eq. (20) $\tilde{L}_i$ denotes the scalar slepton weak doublet, $H_{1,2}$ are the scalar Higgs doublet fields. In the above equations the trilinear terms proportional to λ, $\bar{\lambda}'$, Λ, Λ' and the bilinear terms violate lepton number, while the trilinear terms proportional to $\bar{\lambda}''$, Λ'' violate baryon number conservation. The coupling constants λ ($\bar{\lambda}''$) are antisymmetric in the first (last) two indices. The bar sign in $\bar{\lambda}'$, $\bar{\lambda}''$ denotes that all the definitions are given in the gauge basis for the quark fields. Later on we will change to the mass basis and drop the bars. Using the freedom in the definition of lepton and Higgs superfields we choose the basis where the vacuum expectation values of all the sneutrino fields vanish: $\langle\tilde{\nu}_i\rangle = 0$.

The Lagrangian terms generated by the trilinear terms of the superpotential in Eq. (19) and involving two down quarks needed for the construction of the pARC operators in (2)–(5) are as follows:

$$\mathcal{L}_\lambda = -\lambda'_{ijk}\tilde{\nu}_{iL}\bar{d}_k P_L d_j - \tfrac{1}{2}\lambda''_{ijk}\tilde{u}^\star_{iR}\bar{d}_j P_L d^c_k + \text{H.c.} \tag{21}$$

where d_j stands for the down quark. It can be easily seen that the interactions in Eq. (21) can generate in second order perturbation theory the only $\Delta B = 1$ effective operator contributing to $B^0(\bar{B}^0) \to \phi K^{*0}(\bar{K}^{*0})$-decay. This is the operator $(\bar{s}P_L b)(\bar{s}P_R s)$ which does not belong to the pARC operators listed in Eqs. (2)–(5). Thus, we conclude that the trilinear $\not{R}_p$-couplings alone cannot resolve the polarization anomaly in $B^0(\bar{B}^0) \to \phi K^{*0}(\bar{K}^{*0})$-decay.

Let us see if the bilinear $\not{R}_p$-terms may help in the solution of this problem. The presence of the bilinear terms leads to terms in the scalar potential which are linear in the sneutrino fields, $\tilde{\nu}_i$. First, this results in $\tilde{\nu} - H^0_{1,2}$ mixing. Also, the linear terms drive the $\tilde{\nu}_i$ fields to nonzero vacuum expectation values $\langle\tilde{\nu}_i\rangle \neq 0$ at the minimum of the scalar potential. At this ground state the MSSM vertices $\tilde{Z}\nu\,\tilde{\nu}$ and $\tilde{W}e\tilde{\nu}$ produce the gaugino-lepton mixing mass terms $\tilde{Z}\nu\langle\tilde{\nu}\rangle$, $\tilde{W}e\langle\tilde{\nu}\rangle$ (with $\tilde{W}$, $\tilde{Z}$ being wino and zino fields). These terms taken along with the lepton-higgsino $\mu_i L_i \tilde{H}_1$ mixing from the superpotential of Eq. (19) form 7×7 neutral fermion and 5×5 charged fermion mass matrices [11]. This leads to a nontrivial neutrino mass matrix and Lepton Flavor Violation in the sector of charged leptons. However, these effects are obviously irrelevant for the generation of the effective 4-quark operators.

The above mentioned effect of sneutrino-Higgs mixing $\tilde{\nu} - H^0_{1,2}$ is different. It corresponds to a nondiagonal mass matrix for the neutral scalars $(H^0_1, H^0_2, \tilde{\nu}_e, \tilde{\nu}_\mu, \tilde{\nu}_\tau)$ in the bilinear part of the $\not{R}_p$ scalar potential [12]. From Eqs. (19) and (20) we write

$$V^{\rm soft}_{\not{R}_p} = (\mu^*\mu_j H_1^\dagger + \tilde{\mu}^2_{2j}H_2 + \tilde{\mu}^2_{1j}H_1^\dagger)\tilde{L}_j + \text{H.c.} \tag{22}$$

Using the minimization condition

$$\tilde{\mu}^{\,2}_{1j} + \mu^*\mu_j + \tilde{\mu}^2_{2j}\tan\beta = 0 \tag{23}$$

in the basis of lepton and Higgs superfields where $\langle\tilde{\nu}_i\rangle = 0$ we can rewrite Eq. (22) in the form

$$V^{\rm soft}_{\not{R}_p} = \tilde{\mu}^2_{2j}(H_2 - \tan\beta H_1^\dagger)\tilde{L}_j + \text{H.c.}, \tag{24}$$

where $\tan\beta = \langle H^0_2\rangle/\langle H^0_1\rangle$. Rotating these fields to the mass eigenstate basis we assume smallness of sneutrino-Higgs mixing characterized by the small ratio $(\tilde{\mu}_{kj}/M_{h_{1,2}})^2$, where $\tilde{\mu}^2_{kj}$ is the $\not{R}_p$ soft parameter from Eq. (20) and $M_{h_{1,2}}$ are the neutral Higgs masses [13]. In the leading order in this small parameter we obtain the following interactions of sneutrinos with down quarks and charged leptons

$$\mathcal{L}_{\tilde{\nu}ll} = \eta_j\left[\frac{m_{d_i}}{M_W}(\bar{d}_i d_i) + \frac{m_{l_i}}{M_W}(\bar{l}_i l_i)\right]\tilde{\nu}_j, \tag{25}$$

with the couplings

$$\eta_j = \frac{g_2}{2}\tilde{\mu}^2_{2j}\frac{\tan\beta}{\sqrt{1+\tan^2\beta}}\left(\frac{\cos\alpha}{M^2_{h_2}} - \frac{\sin\alpha}{M^2_{h_1}}\right). \tag{26}$$

Here α is the mixing angle of the neutral Higgses in the limit of no mixing with the sneutrino fields

$$H^0_1 = -\sin\alpha\cdot h^0_1 + \cos\alpha\cdot h^0_2,\qquad H^0_2 = \cos\alpha\cdot h^0_1 + \sin\alpha\cdot h^0_2, \tag{27}$$

where $h^0_{1,2}$ are the corresponding mass eigenstates with the masses M_{h_1}, M_{h_2}. Note that H^0_2, which has no couplings to the down quarks and leptons, does not contribute to Eq. (25).

Now, combining the trilinear and bilinear $\not{R}_p$-interactions from Eq. (21) and (25), as shown in Fig. 1, we obtain in second order perturbation theory the following effective Hamiltonian after integrating out the heavy sneutrino fields

$$\mathcal{H}_{\not{R}_p} = \frac{m_{d_j}}{M_W}(\bar{d}_j d_j)\left(\frac{\eta_i}{m^2_{\tilde{\nu}_i}}\lambda'^{*}_{im3}\bar{d}_m P_R b + \frac{\eta^*_i}{m^2_{\tilde{\nu}_i}}\lambda'_{i3m}\bar{d}_m P_L b\right) + \frac{m_{l_j}}{M_W}(\bar{l}_j l_j)\left(\frac{\eta_i}{m^2_{\tilde{\nu}_i}}\lambda'^{*}_{im3}\bar{d}_m P_R b + \frac{\eta^*_i}{m^2_{\tilde{\nu}_i}}\lambda'_{i3m}\bar{d}_m P_L b\right) + \text{H.c.} \tag{28}$$

The 4-quark terms involve the pARC operators O_{15} and O_{17} from the list of Eqs. (2)–(5) with the following Wilson coefficients:

$$c_{15} = \frac{\sqrt{2}}{G_F}\frac{m_s}{M_W}\frac{\eta_i}{m^2_{\tilde{\nu}_i}}\lambda'^{*}_{i23}, \qquad c_{17} = \frac{\sqrt{2}}{G_F}\frac{m_s}{M_W}\frac{\eta^*_i}{m^2_{\tilde{\nu}_i}}\lambda'_{i32}. \tag{29}$$

Thus $\not{R}_p$ SUSY seems to satisfy the pARC as it allows appropriate operator structures. In the following we have to check if the existing experimental constraints on the $\not{R}_p$-parameters entering into the definition of the Wilson coefficients allow one to accommodate the values of Eqs. (17) and (18).

IV. EXPERIMENTAL CONSTRAINTS ON WILSON COEFFICIENTS

Examining Eq. (28) we note that the strength of both the 4-quark and quark-lepton operators is determined by the same combination of the R-parity conserving and $\not{R}_p$-parameters forming the Wilson coefficients $c_{15,17}$. Therefore, one can directly constrain the $c_{15,17}$ parameters from the existing stringent experimental upper bound on the $B_s \to \mu^+\mu^-$ branching ratio [14]

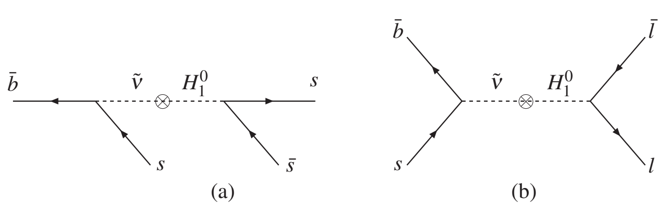

FIG. 1. The $\not{R}_p$ SUSY contribution to the $\bar{b} \to s s \bar{s}$ (a) and to the $\bar{b}s \to l\bar{l}$ transition operators. The sign ⊗ denotes $\not{R}_p$ soft sneutrino-Higgs mixing. The left hand vertices in both diagrams are due to the $\not{R}_p$ superpotential λ' coupling, while the right hand ones correspond to the R-parity conserving $H_1 - q - \bar{q}$ and $H_1 - l - \bar{l}$ Yukawa couplings.

$$\text{Br}(B_s \to \mu^+\mu^-) \leq 1.0 \times 10^{-7}(90\%\text{C.L.}). \tag{30}$$

An important advantage of this constraint is that it applies to the coefficients $c_{15,17}$ as a whole, avoiding uncertainties related to the presence of several R-parity conserving ($\tan\beta$, α, M_{h_1,h_2}, $m_{\tilde{\nu}}$) and violating parameters ($\tilde{\mu}^2_{2j}$, λ').

The contribution of the quark-lepton interactions in the Lagrangian (28) to the decay rate of this process can be written in terms of the Wilson coefficients $c_{15,17}$ as

$$\Gamma(B_s \to \mu^+\mu^-) = \frac{G_F^2}{2}\frac{m_{B_s}}{32\pi}\left(\frac{m_\mu}{m_s}\right)^2\left(f_{B_s}\frac{m^2_{B_s}}{m_b + m_s}\right)^2 \times (c_{15} - c_{17})^2\left[1 - \left(\frac{2m_\mu}{m_{B_s}}\right)^2\right]^{3/2} \tag{31}$$

where we used

$$\langle 0|\bar{s}\gamma_5 b|\bar{B}^0_s\rangle = i f_{B_s}\frac{m^2_{B_s}}{m_b + m_s}. \tag{32}$$

We use the following numerical values for the quantities in above equations: $f_{B_s} = 0.2$ GeV [15], $m_{B_s} = 5.367$ GeV, $m_b = 4.6$ GeV, $m_s = 0.15$ GeV and $\tau_{B_s} = 1.46 \times 10^{-12}$ s [1]. Considering the two scenarios of Ref. [9] as displayed in Eqs. (17) and (18) [denoted as (i) and (ii)] we get from the experimental limit (30) the following upper bounds

$$|c_{15}|, |c_{17}| \leq 1.4 \times 10^{-4}. \tag{33}$$

Using the definitions of Eqs. (15) and (16) these limits can be translated to upper limits on the effective coefficients

$$|\tilde{a}_{23}|, |\tilde{a}_{25}| \leq 5.9 \times 10^{-6}. \tag{34}$$

These limits are about 2 orders of magnitude smaller then the values given in Eqs. (17) and (18) required for the solution of the polarization anomaly.

Thus, we conclude that the polarization anomaly observed in $B^0(\bar{B}^0) \to \phi K^{*0}(\bar{K}^{*0})$ decay by the *BABAR* [2] and BELLE [3] collaborations cannot be explained within the $\not{R}_p$ SUSY framework, despite the occurrence of effective operators with the chiral structure required qualitatively.

As a byproduct of our analysis the limits of Eq. (33) set new upper limits on the products of the soft and superpotential $\not{R}_p$-parameters of Eqs. (26) and (29). Since the expressions for the Wilson coefficients $c_{15,17}$ contain the R-parity conserving parameters as well we choose one representative point in the SUSY parameter space in order to illustrate the limits on the $\not{R}_p$-parameters. We take a typical mSUGRA: the so-called SPS 1a point from the list of nine Snowmass benchmark points [16]. This choice corresponds to $\tan\beta = 10$, $m_0 = -A_0 = 0.25 m_{1/2} = 100$ GeV and $\mu > 0$. For this parameters we find

$$\left(\frac{\tilde{\mu}_{2i}}{100\ \mathrm{GeV}}\right)^2 |\lambda'_{i23}|, \qquad \left(\frac{\tilde{\mu}_{2i}}{100\ \mathrm{GeV}}\right)^2 |\lambda'_{i32}| \leq 5.6 \times 10^{-3}. \tag{35}$$

To our knowledge in the literature (for a review see, for instance [17]) there have not been established experimental limits on these products of $\not{R}_p$-parameters. However, there exist bounds on $\tilde{\mu}^2_{2i}$, λ'_{i23} and λ'_{i32} separately from various low energy processes [17]. This allows one to obtain indirect bounds on their products and compare them with those in Eq. (35). The soft $\not{R}_p$-parameter $\tilde{\mu}^2_{2i}$, contributes to the neutrino mass matrix at one-loop level. Thus it is constrained by the present limits on neutrino masses and mixing from neutrino oscillations. With the SPS 1a set of the *R*-parity conserving parameters one has: $(\tilde{\mu}_{2i}/100\ \mathrm{GeV})^2 \leq 10^{-4}$. Existing constraints on the trilinear $\not{R}_p$-couplings are typically as follows: $\lambda'_{i23}, \lambda'_{i32} \leq 0.2$. Combining these constraints we have the limits

$$\left(\frac{\tilde{\mu}_{2i}}{100\ \mathrm{GeV}}\right)^2 |\lambda'_{i23}|, \qquad \left(\frac{\tilde{\mu}_{2i}}{100\ \mathrm{GeV}}\right)^2 |\lambda'_{i32}| \leq 2.0 \times 10^{-5}. \tag{36}$$

which are 2 orders of magnitude better then those in Eq. (35). Nevertheless, the latter can still be useful as direct constraints on the specific products of the bilinear and trilinear $\not{R}_p$-parameters. Note that these constraints correspond to a particular point in the MSSM parameter space and in some other points the above limits may significantly change. The detailed study of this question is beyond the scope of the present paper.

V. CONCLUSIONS

We analyzed the $\not{R}_p$ SUSY model with respect to its ability to account for the polarization anomaly in $B^0(\bar{B}^0) \to \phi K^{*0}(\bar{K}^{*0})$-decay observed by the *BABAR* [2] and BELLE [3] collaborations. Within this framework we have determined the effective $\Delta B = 1$ operators with chirality structures appropriate for a possible resolution of this anomaly. However, the experimental data on $B \to \mu^+\mu^-$-decay set stringent limits on the respective Wilson coefficients, which are about 2 orders of magnitude below the values required to resolve the polarization anomaly. This gap of 2 orders of magnitude can hardly by eliminated by the uncertainties in the hadronic parameters involved in the calculation of the helicity amplitudes of $B^0(\bar{B}^0) \to \phi K^{*0}(\bar{K}^{*0})$-decay. Therefore, we do not believe that $\not{R}_p$ SUSY is able to account for the *B*-decay polarization anomaly.

As a byproduct we used the experimental data on $B \to \mu^+\mu^-$-decay to set a new upper limit on the product of the two $\not{R}_p$-parameters $\tilde{\mu}^2_{2i}|\lambda'_{i23}|$ and $\tilde{\mu}^2_{2i}|\lambda'_{i32}|$, where $\tilde{\mu}^2_{2i}$ and λ'_{ijk} are bilinear soft and trilinear superpotential $\not{R}_p$-parameters, respectively.

ACKNOWLEDGMENTS

This work was supported in part by CONICYT (Chile) under grants FONDECYT No. 1030244 and PBCT/No. 285/2006 and by the DFG under contracts No. FA67/31-1 and GRK683. This research is also a part of the EU Integrated Infrastructure Initiative Hadronphysics project under the contract No. RII3-CT-2004-506078 and President grant of Russia "Scientific Schools" No. 5103.2006.2. S.K. would like to thank Deutscher Akademischer Austausch Dienst (DAAD) for the support within the program "Förderung ausländischer Gastdozenten zu Lehrtätigkeiten an deutschen Hochschulen". S.K. is also thankful to the Tübingen theory group for hospitality. J.C.H. thanks High Energy Physics LatinAmerican-European Network (HELEN) for the support within the Advanced Training program.

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