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Forward-backward asymmetry, branching ratio and rate difference between electron and muon channels of $B \to K_1(K^*)\ell^+\ell^-$ transition in supersymmetric models

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Abstract: The mass eigen states $K_1(1270)$ and $K_1(1400)$ are mixture of the strange members of two axial-vector SU(3) octet, $^3P_1(K_1^A)$ and $^1P_1(K_1^B)$. Taking into account this mixture, the forward-backward asymmetry(A_{FB}), branching ratio(Br) and rate difference of electron channel to muon channel(R) of $B \to K_1(1270, 1400)\ell^+\ell^-$ transitions are studied in the framework of different supersymmetric models. MSSM with R parity is considered because considerable deviation from the standard model predictions can be obtained in $B \to X_s\ell^-\ell^+$. Taking C_{Q1} and C_{Q2} about one which is consistent with the $B \to K^*\mu^+\mu^-$ rate at low dileptonic invariant mass region($1 \leq q^2 \leq 6\mathrm{GeV}^2$), we obtain a size able deviation for A_{FB}, Br and R with respect to the Standard Model results. Any measurement of physical observables and their comparison with the results obtained in this paper can gives useful information about the nature of interactions beyond the standard model.

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Contents

1 Introduction

The Standard Model (SM) explains all experimental predictions well. Despite all the success of SM, we can not accept that it is the ultimate theory of nature since there are many questions to be discussed. Some issues such as gauge and fermion mass hierarchy, matter- antimatter asymmetry, number of generations, the nature of the dark matter and the unification of fundamental forces can not be addressed by the SM. In other words, the SM can be considered as an effective theory of some fundamental theory at low energy.

One of the most reasonable extension of the SM is the Supersymmetry (SUSY) [1]. It is an important element in the string theory, which is the most-favored candidate for unifying the all known interactions including gravity. The SUSY is assumed to contribute to overcome the mass hierarchy problem between m_W and the Planck scale via canceling the quadratic divergences in the radiative corrections to the Higgs boson mass-squared [2].

To verify the SUSY theories, we need to explore the supersymmetric particles (sparticles). Two types of studies can be conducted to examine these sparticles. In the direct search, the center of mass energy of colliding particles should be increased to produce SUSY particles at the TeV scale, hence, it will be accessible to the LHC. On the other hand, we can look for SUSY effects, indirectly. The sparticles can contribute to the transitions at loop level. The flavor changing neutral current (FCNC) transition of $b \to s$ induced by quantum loop level can be considered as a good candidate for studying the possible effects of sparticles. For the most recent studies in this regard see ref. [3] and the references therein.

The $B \to K_1 \ell^+ \ell^-$ transition proceeds via the FCNC transition of $b \to s$ at quark level. $b \to s$ transition is the most sensitive and stringiest test for the SM at one loop level, where, it is forbidden in SM at tree level [4, 5]. Although, the FCNC transitions have small branching fractions, quite intriguing results are obtained in ongoing experiments. The inclusive $B \to X_s \ell^+ \ell^-$ decay is observed in BaBaR [6] and Belle collaborations. These collaborations have also announced the measuring exclusive modes $B \to K \ell^+ \ell^-$ [7–9] and $B \to K^* \ell^+ \ell^-$ [10]. The obtained experimental results on these transitions are in a good consistency with theoretical predictions [11–19] the results of which can be used to constrain the new physics (NP) effects.

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In the present work, calculating the forward-backward asymmetry and the branching fraction, we investigate the possible effects of supersymmetric theories on the branching ratio of $B \to K_1\ell^+\ell^-$ transition. Experimentally, the $K_1(1270)$ and $K_1(1400)$ are the mixtures of the strange members of the two axial-vector SU(3) octet ${}^3P_1(K_1^A)$ and ${}^1P_1(K_1^B)$. The $K_1(1270, 1400)$ and $K_1^{A,B}$ states are related to each other as [20]

$$\begin{pmatrix} |\overline{K}_1(1270)\rangle \\ |\overline{K}_1(1400)\rangle \end{pmatrix} = M \begin{pmatrix} |\overline{K}_{1A}\rangle \\ |\overline{K}_{1B}\rangle \end{pmatrix}, \quad \text{with} \quad M = \begin{pmatrix} \sin\theta_{K_1} & \cos\theta_{K_1} \\ \cos\theta_{K_1} & -\sin\theta_{K_1} \end{pmatrix}. \tag{1.1}$$

The branching ratio of the $K_1(1400)$ case is smaller than the $K_1(1270)$ [20], so we consider only $B \to K_1(1270)\ell^+\ell^-$. Note that lepton polarization and angular distribution of this decay in the frame work of SM has recently been studied in refs. [21, 22].

The radiative B decay involving the $K_1(1270)$, the orbitally excited (P-wave) state, is observed by BELLE [23] and other radiative and semileptonic decay modes involving $K_1(1270)$ and $K_1(1400)$ are hopefully expected to be seen soon. Just like $B \to K^*(892)\ell^+\ell^-$ decays, $B \to K_1\ell^+\ell^-$ decays can offer the good probe to the new physics effects, and are much more sophisticated due to the mixing of the K_{1A} and K_{1B}, which are the 1^3P_1 and 1^1P_1 states, respectively [20].

The outline of the paper is as follows: In section 2, we calculate the decay amplitude and forward-backward asymmetry of the $B \to K_1\ell^+\ell^-$ transition within SUSY models. Section 3 is devoted to the numerical analysis and discussion of the considered transition as well as our conclusions.

2 The effective Hamiltonian

In the most SUSY models R parity is conserved so that SUSY contributions on a physical observable appear at the quantum loop level. The QCD corrected effective Lagrangian for the decays $b \to s(d)\ell^+\ell^-$ can be achieved by integrating out the heavy quarks and the heavy electroweak bosons in the SUSY models with R parity:

$$\begin{aligned} \mathcal{H}_{\text{eff}} = \frac{G_F\alpha V_{tb}V_{ts}^*}{2\sqrt{2}\pi}\Big[& C_9^{\text{eff}}(m_b)\bar{s}\gamma_\mu(1-\gamma_5)b\,\bar{\ell}\gamma^\mu\ell + C_{10}(m_b)\bar{s}\gamma_\mu(1-\gamma_5)b\,\bar{\ell}\gamma^\mu\gamma_5\ell \\ & -2m_bC_7(m_b)\frac{1}{q^2}\bar{s}i\sigma_{\mu\nu}q^\nu(1+\gamma_5)b\,\bar{\ell}\gamma^\mu\ell + C_{Q_1}\bar{s}(1+\gamma_5)b\,\bar{\ell}\ell + C_{Q_2}\bar{s}(1+\gamma_5)b\,\bar{\ell}\gamma_5\ell\Big], \end{aligned} \tag{2.1}$$

SUSY introduces several additional classes of contributions including; I. gluino, down-type squark loop, II. chargino, up-type squark loop, III. chargino, up-type squark loop, (Higgs field attaching to charginos) and IV. neutralino down-type squark loop [24]. The neutral Higgs couplings SUSY contributions are mainly involved via terms proportional with $C_{Q_{1,2}}$. These additional terms with respect to the SM come from the neutral Higgs bosons(NHBs) exchange diagrams, whose manifest forms and corresponding Wilson coefficients can be found in [25–31]: The coefficients $C_{Qi}(m_W)$ in the MSSM with R parity and the results are [28, 29]:

$$
\begin{aligned}
C_{Q1}(m_W) =& \frac{m_b m_\ell}{4m_{h^0}^2\sin^2\theta_W}\mathrm{tg}^2\beta\Bigg\{(\sin^2\alpha+h\cos^2\alpha)\Bigg[\frac{1}{x_{Wt}}(f_1(x_{Ht})-f_1(x_{Wt}))\\
&+\sqrt{2}\sum_{i=1}^{2}\frac{m_{\chi_i}}{m_W}\frac{U_{i2}}{\cos\beta}\left(-V_{i1}f_1(x_{\chi_i\tilde{q}})+\sum_{k=1}^{2}\Lambda(i,k)T_{k1}f_1(x_{\chi_i\tilde{t}_k})\right)\\
&+\left(1+\frac{m_{H_\pm}^2}{m_W^2}\right)f_2(x_{Ht},x_{Wt})\Bigg]-\frac{m_{h_0}^2}{m_W^2}f_2(x_{Ht},x_{Wt})\\
&+2\sum_{ii'=1}^{2}(B_1(i,i')\Gamma_1(i,i')+A_1(i,i')\Gamma_2(i,i'))\Bigg\}\\
C_{Q2}(m_W) =& -\frac{m_b m_\ell}{4m_{A^0}^2\sin^2\theta_W}\mathrm{tg}^2\beta\Bigg\{\frac{1}{x_{Wt}}(f_1(x_{Ht})-f_1(x_{Wt}))+2f_2(x_{Ht},x_{Wt})\\
&+\sqrt{2}\sum_{i=1}^{2}\frac{m_{\chi_i}}{m_W}\frac{U_{i2}}{\cos\beta}\left(-V_{i1}f_1(x_{\chi_i\tilde{q}})+\sum_{k=1}^{2}\Lambda(i,k)T_{k1}f_1(x_{\chi_i\tilde{t}_k})\right)\\
&+2\sum_{ii'=1}^{2}(-U_{i'2}V_{i1}\Gamma_1(i,i')+U_{i2}^*V_{i'1}^*\Gamma_2(i,i'))\Bigg\}
\end{aligned}
\tag{2.2}
$$

where

$$
\begin{aligned}
B_1(i,i') =& \left(-\frac{1}{2}U_{i'1}V_{i2}\sin 2\alpha(1-h)+U_{i'2}V_{i1}(\sin^2\alpha+h\cos^2\alpha)\right)\\
A_1(i,i') =& \left(-\frac{1}{2}U_{i1}^*V_{i'2}^*\sin 2\alpha(1-h)+U_{i2}^*V_{i'1}^*(\sin^2\alpha+h\cos^2\alpha)\right)\\
\Gamma_1(i,i') =& m_{\chi_i}m_{\chi_{i'}}U_{i2}\left(-\frac{1}{\tilde{m}^2}f_2(x_{\chi_i\tilde{q}},x_{\chi_{i'}\tilde{q}})V_{i'1}+\sum_{k=1}^{2}\frac{1}{m_{\tilde{t}_k}^2}\Lambda(i',k)T_{k1}f_2(x_{\chi_i\tilde{t}_k},x_{\chi_{i'}\tilde{t}_k})\right)\\
\Gamma_2(i,i') =& U_{i2}\left(-f_2(x_{\chi_i\tilde{q}},x_{\chi i'\tilde{q}})V_{i'1}+\sum_{k=1}^{2}\Lambda(i',k)T_{k1}f_2(x_{\chi_i\tilde{t}_k},x_{\chi_{i'}\tilde{t}_k})\right)\\
\Lambda(i,k) =& V_{i1}T_{k1}-V_{i2}T_{k2}\frac{m_t}{\sqrt{2}m_W\sin\beta}\\
f_1(x_{ij}) =& 1-\frac{x_{ij}}{x_{ij}-1}ln\ x_{ij}+ln\ x_{Wj}\\
f_2(x,y) =& \frac{1}{x-y}\left(\frac{x}{x-1}ln\ x-\frac{y}{y-1}ln\ y\right)\\
x_{ij} =& m_i^2/m_j^2
\end{aligned}
\tag{2.3}
$$

with m_i being the mass of the particle i, and

$$
\begin{aligned}
C_{Q3}(m_w) =& \frac{m_b e^2}{m_\ell g^2}\{C_{Q1}(m_w)+C_{Q2}(m_w)\}\\
C_{Q4}(m_w) =& \frac{m_b e^2}{m_\ell g^2}\{C_{Q1}(m_w)-C_{Q2}(m_w)\}\\
C_{Qi}(m_w) =& 0,(i=5,\cdots 10)
\end{aligned}
\tag{2.4}
$$

In eqs. (2.2) and (2.3), U and V are matrices which diagonalize the mass matrix of charginos, T is the matrix reflecting the mixing of stops t_R and t_L, m_{χ_j} denote the chargino masses, $\tilde{m}$ is the average mass of u-type squarks $\tilde{q}$ of the first two generations, h is the square of the ratio of the mass of h^0 to the mass of H^0 and α is the mixing angle of neutral components of the two higgs doublets in the model. And in eq. (2.2) less important terms have been omitted because they are numerically negligible compared to those given in eq. (2.2) when tan$\beta \geq 20$.

The effects of new scalar and pseudoscalar type interactions on physical observables come through the terms which are proportional to the mass of final state leptons. The effects of the other contributions come through the modification of known SM Wilson coefficients. The C_i in the frame work of SM are calculated in naive dimensional regularization (NDR) scheme at the leading order(LO), next to leading order(NLO) and next-to-next leading order (NNLO) in the SM [32]–[39]. $C_9^{\rm eff}(\hat{s}) = C_9 + Y(\hat{s})$, where $Y(\hat{s}) = Y_{\rm pert}(\hat{s}) + Y_{\rm LD}$ contains both the perturbative part $Y_{\rm pert}(\hat{s})$ and long-distance part $Y_{\rm LD}(\hat{s})$. $Y(\hat{s})_{\rm pert}$ is given by [32]

$$\begin{aligned} Y_{\rm pert}(\hat{s}) &= g(\hat{m}_c, \hat{s})c_0 \\ &-\frac{1}{2}g(1,\hat{s})(4\bar{c}_3 + 4\bar{c}_4 + 3\bar{c}_5 + \bar{c}_6) - \frac{1}{2}g(0,\hat{s})(\bar{c}_3 + 3\bar{c}_4) \\ &+\frac{2}{9}(3\bar{c}_3 + \bar{c}_4 + 3\bar{c}_5 + \bar{c}_6), \end{aligned} \tag{2.5}$$

$$\text{with} \quad c_0 \equiv \bar{c}_1 + 3\bar{c}_2 + 3\bar{c}_3 + \bar{c}_4 + 3\bar{c}_5 + \bar{c}_6, \tag{2.6}$$

and the function $g(x, y)$ is defined in [32]. Here, $\bar{c}_1$ — $\bar{c}_6$ are the Wilson coefficients in the leading logarithmic approximation. The relevant Wilson coefficients are given in refs. [40]. $Y(\hat{s})_{\rm LD}$ involves $B \to K_1 V(\bar{c}c)$ resonances [33], where $V(\bar{c}c)$ are the vector charmonium states. We follow refs. [33, 41] and set

$$Y_{\rm LD}(\hat{s}) = -\frac{3\pi}{\alpha_{\rm em}^2}c_0 \sum_{V=\psi(1s),\cdots} \kappa_V \frac{\hat{m}_V \mathcal{B}(V \to \ell^+\ell^-)\hat{\Gamma}_{\rm tot}^V}{\hat{s} - \hat{m}_V^2 + i\hat{m}_V\hat{\Gamma}_{\rm tot}^V}, \tag{2.7}$$

where $\hat{\Gamma}_{\rm tot}^V \equiv \Gamma_{\rm tot}^V/m_B$ and κ_V takes different value for different exclusive semileptonic decay. The relevant properties of vector charmonium states are summarized in table 1. The Wilson coefficients in the frame work of the SUSY can be different from the their SM values. While the SUSY effects on C_7, which is proportional to the product of the top and bottom Yukawa coupling constant, $m_t m_b \tan\beta/\sin^2\beta$, is sizable for large $\tan\beta$, there are no such effects in the calculation of C_9 and C_{10}.

One has to sandwich the inclusive effective Hamiltonian between initial hadron state $B(p_B)$ and final hadron state K_1 in order to obtain the matrix element for the exclusive decay $B \to K_1\ell^+\ell^-$. Following from eq. (2.1), in order to calculate the decay width and other physical observable of the exclusive $B \to K_1\ell^+\ell^-$ decay, we need to parameterize the matrix elements in terms of formfactors.

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V	Mass[GeV]	$\Gamma^V_{\rm tot}$[MeV]	$\mathcal{B}(V \to \ell^+\ell^-)$	
$J/\Psi(1S)$	3.097	0.093	5.9×10^{-2}	for $\ell = e, \mu$
$\Psi(2S)$	3.686	0.327	7.4×10^{-3}	for $\ell = e, \mu$
			3.0×10^{-3}	for $\ell = \tau$
$\Psi(3770)$	3.772	25.2	9.8×10^{-6}	for $\ell = e$
$\Psi(4040)$	4.040	80	1.1×10^{-5}	for $\ell = e$
$\Psi(4160)$	4.153	103	8.1×10^{-6}	for $\ell = e$
$\Psi(4415)$	4.421	62	9.4×10^{-6}	for $\ell = e$

Table 1. Masses, total decay widths and branching fractions of dilepton decays of vector charmonium states [42].

The $\overline{B}(p_B) \to \overline{K}_1(p_{K_1}, \lambda)$ form factors are defined by [20]

$$
\begin{aligned}
&\langle \overline{K}_1(p_{K_1}, \lambda)|\bar{s}\gamma_\mu(1-\gamma_5)b|\overline{B}(p_B)\rangle \\
&\quad = -i\frac{2}{m_B + m_{K_1}}\epsilon_{\mu\nu\rho\sigma}\varepsilon^{*\nu}_{(\lambda)}p^\rho_B p^\sigma_{K_1} A^{K_1}(q^2) \\
&\qquad - \left[(m_B + m_{K_1})\varepsilon^{(\lambda)*}_\mu V^{K_1}_1(q^2) - (p_B + p_{K_1})_\mu(\varepsilon^*_{(\lambda)} \cdot p_B)\frac{V^{K_1}_2(q^2)}{m_B + m_{K_1}}\right] \\
&\qquad + 2m_{K_1}\frac{\varepsilon^*_{(\lambda)} \cdot p_B}{q^2}q_\mu\left[V^{K_1}_3(q^2) - V^{K_1}_0(q^2)\right], \qquad (2.8)\\
&\langle \overline{K}_1(p_{K_1}, \lambda)|\bar{s}\sigma_{\mu\nu}q^\nu(1+\gamma_5)b|\overline{B}(p_B)\rangle \\
&\quad = 2T^{K_1}_1(q^2)\epsilon_{\mu\nu\rho\sigma}\varepsilon^{*\nu}_{(\lambda)}p^\rho_B p^\sigma_{K_1} \\
&\qquad - iT^{K_1}_2(q^2)\left[(m^2_B - m^2_{K_1})\varepsilon^{(\lambda)}_{*\mu} - (\varepsilon^*_{(\lambda)} \cdot q)(p_B + p_{K_1})_\mu\right] \\
&\qquad - iT^{K_1}_3(q^2)(\varepsilon^*_{(\lambda)} \cdot q)\left[q_\mu - \frac{q^2}{m^2_B - m^2_{K_1}}(p_{K_1} + p_B)_\mu\right], \qquad (2.9)
\end{aligned}
$$

where $q \equiv p_B - p_{K_1} = p_{\ell^+} + p_{\ell^-}$. By multiplying both sides of eq. (2.8) with q^μ, one can obtain the expression in terms of form factors for $\langle \overline{K}_1(p_{K_1}, \lambda)|\bar{s}(1+\gamma_5)b|\overline{B}(p_B)\rangle$

$$
\begin{aligned}
&\langle \overline{K}_1(p_{K_1}, \lambda)|\bar{s}(1+\gamma_5)b|\overline{B}(p_B)\rangle \\
&\quad = \frac{1}{m_b + m_s}\Big\{ \\
&\qquad - \left[(m_B + m_{K_1})(\varepsilon^{(\lambda)*}.q)V^{K_1}_1(q^2) - (m_B - m_{k_1})(\varepsilon^*_{(\lambda)} \cdot p_B)V^{K_1}_2(q^2)\right] \\
&\qquad + 2m_{K_1}(\varepsilon^*_{(\lambda)} \cdot p_B)\left[V^{K_1}_3(q^2) - V^{K_1}_0(q^2)\right]\Big\}, \qquad (2.10)
\end{aligned}
$$

The formfactors of $B \to K_1(1270)$ and $B \to K_1(1400)$ can be expressed in terms of

$B \to K_A$ and $B \to K_B$ as follows(see [20]):

$$\begin{pmatrix} \langle \overline{K}_1(1270)|\bar{s}\gamma_\mu(1-\gamma_5)b|\overline{B}\rangle \\ \langle \overline{K}_1(1400)|\bar{s}\gamma_\mu(1-\gamma_5)b|\overline{B}\rangle \end{pmatrix} = M \begin{pmatrix} \langle \overline{K}_{1A}|\bar{s}\gamma_\mu(1-\gamma_5)b|\overline{B}\rangle \\ \langle \overline{K}_{1B}|\bar{s}\gamma_\mu(1-\gamma_5)b|\overline{B}\rangle \end{pmatrix}, \tag{2.11}$$

$$\begin{pmatrix} \langle \overline{K}_1(1270)|\bar{s}\sigma_{\mu\nu}q^\nu(1+\gamma_5)b|\overline{B}\rangle \\ \langle \overline{K}_1(1400)|\bar{s}\sigma_{\mu\nu}q^\nu(1+\gamma_5)b|\overline{B}\rangle \end{pmatrix} = M \begin{pmatrix} \langle \overline{K}_{1A}|\bar{s}\sigma_{\mu\nu}q^\nu(1+\gamma_5)b|\overline{B}\rangle \\ \langle \overline{K}_{1B}|\bar{s}\gamma_{\mu\nu}q^\nu(1+\gamma_5)b|\overline{B}\rangle \end{pmatrix}, \tag{2.12}$$

using the mixing matrix M being given in eq. (1.1) the formfactors $A^{K_1}, V^{K_1}_{0,1,2}$ and $T^{K_1}_{1,2,3}$ can be written as follows:

$$\begin{pmatrix} A^{K_1(1270)}/(m_B + m_{K_1(1270)}) \\ A^{K_1(1400)}/(m_B + m_{K_1(1400)}) \end{pmatrix} = M \begin{pmatrix} A^{K_{1A}}/(m_B + m_{K_{1A}}) \\ A^{K_{1B}}/(m_B + m_{K_{1B}}) \end{pmatrix}, \tag{2.13}$$

$$\begin{pmatrix} (m_B + m_{K_1(1270)})V_1^{K_1(1270)} \\ (m_B + m_{K_1(1400)})V_1^{K_1(1400)} \end{pmatrix} = M \begin{pmatrix} (m_B + m_{K_{1A}})V_1^{K_{1A}} \\ (m_B + m_{K_{1B}})V_1^{K_{1B}} \end{pmatrix}, \tag{2.14}$$

$$\begin{pmatrix} V_2^{K_1(1270)}/(m_B + m_{K_1(1270)}) \\ V_2^{K_1(1400)}/(m_B + m_{K_1(1400)}) \end{pmatrix} = M \begin{pmatrix} V_2^{K_{1A}}/(m_B + m_{K_{1A}}) \\ V_2^{K_{1B}}/(m_B + m_{K_{1B}}) \end{pmatrix}, \tag{2.15}$$

$$\begin{pmatrix} m_{K_1(1270)}V_0^{K_1(1270)} \\ m_{K_1(1400)}V_0^{K_1(1400)} \end{pmatrix} = M \begin{pmatrix} m_{K_{1A}}V_0^{K_{1A}} \\ m_{K_{1B}}V_0^{K_{1B}} \end{pmatrix}, \tag{2.16}$$

$$\begin{pmatrix} T_1^{K_1(1270)} \\ T_1^{K_1(1400)} \end{pmatrix} = M \begin{pmatrix} T_1^{K_{1A}} \\ T_1^{K_{1B}} \end{pmatrix}, \tag{2.17}$$

$$\begin{pmatrix} (m_B^2 - m_{K_1(1270)}^2)T_2^{K_1(1270)} \\ (m_B^2 - m_{K_1(1400)}^2)T_2^{K_1(1400)} \end{pmatrix} = M \begin{pmatrix} (m_B^2 - m_{K_{1A}}^2)T_2^{K_{1A}} \\ (m_B^2 - m_{K_{1B}}^2)T_2^{K_{1B}} \end{pmatrix}, \tag{2.18}$$

$$\begin{pmatrix} T_3^{K_1(1270)} \\ T_3^{K_1(1400)} \end{pmatrix} = M \begin{pmatrix} T_3^{K_{1A}} \\ T_3^{K_{1B}} \end{pmatrix}, \tag{2.19}$$

where it is supposed that $p^\mu_{K_1(1270),K_1(1400)} \simeq p^\mu_{K_{1A}} \simeq p^\mu_{K_{1B}}$ [20]. These formfactors within light-cone QCD sum rule (LCQSR) are estimated in [43].

Thus the matrix element for $B \to K_1\ell^+\ell^-$ in terms of formfactor is given by

$$\begin{aligned}\mathcal{M} = {} & \frac{G_F\alpha_{\rm em}}{2\sqrt{2}\pi}V_{ts}^* V_{tb}\, m_B \cdot (-i) \\ & \left\{\mathcal{T}_\mu^{(K_1),1}\bar{\ell}\gamma^\mu\ell + \mathcal{T}_\mu^{(K_1),2}\bar{\ell}\gamma^\mu\gamma_5\ell + \mathcal{T}^{(K_1),3}\bar{\ell}\ell + \mathcal{T}^{(K_1),4}\bar{\ell}\gamma_5\ell\right\},\end{aligned} \tag{2.20}$$

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where

$$
\begin{aligned}
\mathcal{T}_\mu^{(K_1),1} &= \mathcal{A}^{K_1}(\hat{s})\epsilon_{\mu\nu\rho\sigma}\varepsilon^{*\nu}\hat{p}_B^\rho\hat{p}_{K_1}^\sigma - i\mathcal{B}^{K_1}(\hat{s})\varepsilon_\mu^* \\
&\quad + i\mathcal{C}^{K_1}(\hat{s})(\varepsilon^*\cdot\hat{p}_B)\hat{p}_\mu + i\mathcal{D}^{K_1}(\hat{s})(\varepsilon^*\cdot\hat{p}_B)\hat{q}_\mu, && (2.21)\\
\mathcal{T}_\mu^{(K_1),2} &= \mathcal{E}^{K_1}(\hat{s})\epsilon_{\mu\nu\rho\sigma}\varepsilon^{*\nu}\hat{p}_B^\rho\hat{p}_{K_1}^\sigma - i\mathcal{F}^{K_1}(\hat{s})\varepsilon_\mu^* \\
&\quad + i\mathcal{G}^{K_1}(\hat{s})(\varepsilon^*\cdot\hat{p}_B)\hat{p}_\mu + i\mathcal{H}^{K_1}(\hat{s})(\varepsilon^*\cdot\hat{p}_B)\hat{q}_\mu, && (2.22)\\
\mathcal{T}^{(K_1),3} &= i\mathcal{I}_1^{K_1}(\hat{s})\frac{(\varepsilon^{(\lambda)*}.\hat{q})}{1+\hat{m}_s} + i\mathcal{J}_1^{K_1}(\hat{s})\frac{(\varepsilon^{(\lambda)*}.\hat{p}_B)}{1+\hat{m}_s} && (2.23)\\
\mathcal{T}^{(K_1),4} &= i\mathcal{I}_2^{K_1}(\hat{s})\frac{(\varepsilon^{(\lambda)*}.\hat{q})}{1+\hat{m}_s} + i\mathcal{J}_2^{K_1}(\hat{s})\frac{(\varepsilon^{(\lambda)*}.\hat{p}_B)}{1+\hat{m}_s} && (2.24)
\end{aligned}
$$

with $\hat{p} = p/m_B$, $\hat{p}_B = p_B/m_B$, $\hat{q} = q/m_B$, $\hat{m}_s = m_s/m_B$,and $p = p_B + p_{K_1}$, $q = p_B - p_{K_1} = p_{\ell^+} + p_{\ell^-}$. Here $\mathcal{A}^{K_1}(\hat{s}), \cdots, \mathcal{H}^{K_1}(\hat{s})$ are defined by

$$
\begin{aligned}
\mathcal{A}^{K_1}(\hat{s}) &= \frac{2}{1+\sqrt{\hat{r}_{K_1}}}c_9^{\text{eff}}(\hat{s})A^{K_1}(\hat{s}) + \frac{4\hat{m}_b}{\hat{s}}c_7^{\text{eff}}T_1^{K_1}(\hat{s}), && (2.25)\\
\mathcal{B}^{K_1}(\hat{s}) &= (1+\sqrt{\hat{r}_{K_1}})\left[c_9^{\text{eff}}(\hat{s})V_1^{K_1}(\hat{s}) + \frac{2\hat{m}_b}{\hat{s}}(1-\sqrt{\hat{r}_{K_1}})c_7^{\text{eff}}T_2^{K_1}(\hat{s})\right], && (2.26)\\
\mathcal{C}^{K_1}(\hat{s}) &= \frac{1}{1-\hat{r}_{K_1}}\left[(1-\sqrt{\hat{r}_{K_1}})c_9^{\text{eff}}(\hat{s})V_2^{K_1}(\hat{s}) + 2\hat{m}_b c_7^{\text{eff}}\left(T_3^{K_1}(\hat{s}) + \frac{1-\sqrt{\hat{r}_{K_1}}^2}{\hat{s}}T_2^{K_1}(\hat{s})\right)\right], && (2.27)\\
\mathcal{D}^{K_1}(\hat{s}) &= \frac{1}{\hat{s}}\Bigg[c_9^{\text{eff}}(\hat{s})\left\{(1+\sqrt{\hat{r}_{K_1}})V_1^{K_1}(\hat{s}) - (1-\sqrt{\hat{r}_{K_1}})V_2^{K_1}(\hat{s}) - 2\sqrt{\hat{r}_{K_1}}V_0^{K_1}(\hat{s})\right\} \\
&\quad -2\hat{m}_b c_7^{\text{eff}}T_3^{K_1}(\hat{s})\Bigg], && (2.28)\\
\mathcal{E}^{K_1}(\hat{s}) &= \frac{2}{1+\sqrt{\hat{r}_{K_1}}}c_{10}A^{K_1}(\hat{s}), && (2.29)\\
\mathcal{F}^{K_1}(\hat{s}) &= (1+\sqrt{\hat{r}_{K_1}})c_{10}V_1^{K_1}(\hat{s}), && (2.30)\\
\mathcal{G}^{K_1}(\hat{s}) &= \frac{1}{1+\sqrt{\hat{r}_{K_1}}}c_{10}V_2^{K_1}(\hat{s}), && (2.31)\\
\mathcal{H}^{K_1}(\hat{s}) &= \frac{1}{\hat{s}}c_{10}\left[(1+\sqrt{\hat{r}_{K_1}})V_1^{K_1}(\hat{s}) - (1-\sqrt{\hat{r}_{K_1}})V_2^{K_1}(\hat{s}) - 2\sqrt{\hat{r}_{K_1}}V_0^{K_1}(\hat{s})\right], && (2.32)\\
\mathcal{I}_1^{K_1}(\hat{s}) &= -C_{Q_1}(1+\sqrt{\hat{r}_{K_1}})V_1^{K_1}(\hat{s}) && (2.33)\\
\mathcal{J}_1^{K_1}(\hat{s}) &= C_{Q_1}\{(1+\sqrt{\hat{r}_{K_1}})V_2^{K_1}(\hat{s}) + 2\sqrt{\hat{r}_{K_1}}[V_3^{K_1}(\hat{s}) - V_0^{K_1}(\hat{s})]\} && (2.34)\\
\mathcal{I}_2^{K_1}(\hat{s}) &= \mathcal{I}_1^{K_1}(\hat{s})(C_{Q_2}\to C_{Q_1}), \quad \mathcal{J}_2^{K_1}(\hat{s}) = \mathcal{J}_1^{K_1}(\hat{s})(C_{Q_2}\to C_{Q_1}) && (2.35)
\end{aligned}
$$

with $\hat{r}_{K_1} = m_{K_1}^2/m_B^2$ and $\hat{s} = q^2/m_B^2$.

The dilepton invariant mass spectrum of the lepton pair for the $\overline{B}\to\overline{K}_1\ell^+\ell^-$ decay is given by

$$
\frac{d\Gamma(\overline{B}\to\overline{K}_1\ell^+\ell^-)}{d\hat{s}} = \frac{G_F^2\alpha_{\text{em}}^2 m_B^5}{2^{12}\pi^5}\left|V_{tb}V_{ts}^*\right|^2 v\sqrt{\lambda}\Delta(\hat{s}) \tag{2.36}
$$

where $v = \sqrt{1-4\hat{m}_\ell^2/\hat{s}}$, $\lambda = 1+\hat{r}_{K_1}^2+\hat{s}^2-2\hat{s}-2\hat{r}_{K_1}(1+\hat{s})$ and

$$\begin{aligned}
\Delta(\hat{s}) =& \frac{8Re[\mathcal{F}\mathcal{H}^*]\hat{m}_\ell^2\lambda}{\hat{r}_{K_1}} + \frac{8Re[\mathcal{G}\mathcal{H}^*]\hat{m}_\ell^2(-1+\hat{r}_{K_1})\lambda}{\hat{r}_{K_1}} - \frac{8|\mathcal{H}|^2\hat{m}_\ell^2\hat{s}\lambda}{\hat{r}_{K_1}} \\
&-\frac{2Re[\mathcal{B}\mathcal{C}^*](-1+\hat{r}_{K_1}+\hat{s})(3+3\hat{r}_{K_1}^2-6\hat{s}+3\hat{s}^2-6\hat{r}_{K_1}(1+\hat{s})-v^2\lambda)}{3\hat{r}_{K_1}} \\
&-\frac{|\mathcal{C}|^2\lambda(3+3\hat{r}_{K_1}^2-6\hat{s}+3\hat{s}^2-6\hat{r}_{K_1}(1+\hat{s})-v^2\lambda)}{3\hat{r}_{K_1}} \\
&-\frac{|\mathcal{G}|^2\lambda(3+3\hat{r}_{K_1}^2+12\hat{m}_\ell^2(2+2\hat{r}_{K_1}-\hat{s})-6\hat{s}+3\hat{s}^2-6\hat{r}_{K_1}(1+\hat{s})-v^2\lambda)}{3\hat{r}_{K_1}} \\
&+\frac{|\mathcal{F}|^2(-3-3\hat{r}_{K_1}^2+6\hat{r}_{K_1}(1+16\hat{m}_\ell^2-3\hat{s})+6\hat{s}-3\hat{s}^2+v^2\lambda)}{3\hat{r}_{K_1}} \\
&+\frac{|\mathcal{B}|^2(-3-3\hat{r}_{K_1}^2+6\hat{s}-3\hat{s}^2-6\hat{r}_{K_1}(-1+8\hat{m}_\ell^2+3\hat{s})+v^2\lambda)}{3\hat{r}_{K_1}} \\
&+\frac{2}{3\hat{r}_{K_1}}Re[\mathcal{F}\mathcal{G}^*](12\hat{m}_\ell^2\lambda-(-1+\hat{r}_{K_1}+\hat{s})(3+3\hat{r}_{K_1}^2-6\hat{s}+3\hat{s}^2-6\hat{r}_{K_1}(1+\hat{s})-v^2\lambda)) \\
&+|\mathcal{A}|^2\left(-4\hat{m}_\ell^2\lambda-\frac{\hat{s}}{3}(3+3\hat{r}_{K_1}^2-6\hat{s}+3\hat{s}^2-6\hat{r}_{K_1}(1+\hat{s})+v^2\lambda)\right) \\
&+|\mathcal{E}|^2\left(4\hat{m}_\ell^2\lambda-\frac{\hat{s}}{3}(3+3\hat{r}_{K_1}^2-6\hat{s}+3\hat{s}^2-6\hat{r}_{K_1}(1+\hat{s})+v^2\lambda)\right) \\
&+\lambda\Bigg\{\frac{(4\hat{m}_\ell^2-\hat{s})\,|\mathcal{I}_1|^2}{\hat{r}_{K_1}}+\frac{|\mathcal{J}_1|^2\,(4\hat{m}_\ell^2-\hat{s})}{\hat{r}_{K_1}}+\frac{2Re[\mathcal{I}_1\mathcal{J}_1^*]\,(4\hat{m}_\ell^2-\hat{s})}{\hat{r}_{K_1}}-\frac{|\mathcal{I}_2|^2\hat{s}}{\hat{r}_{K_1}} \\
&-\frac{|\mathcal{J}_2|^2\hat{s}}{\hat{r}_{K_1}}-\frac{2Re[\mathcal{I}_1\mathcal{J}_1^*]\hat{s}}{\hat{r}_{K_1}}+\frac{4Re[\mathcal{H}\mathcal{I}_2^*]\hat{m}_\ell\hat{s}}{\hat{r}_{K_1}}+\frac{4Re[\mathcal{H}\mathcal{J}_2^*]\hat{m}_\ell\hat{s}}{\hat{r}_{K_1}}-\frac{4Re[\mathcal{F}\mathcal{I}_2^*]\hat{m}_\ell}{\hat{r}_{K_1}} \\
&-\frac{4Re[\mathcal{F}\mathcal{J}_2^*]\hat{m}_\ell}{\hat{r}_{K_1}}-\frac{4Re[\mathcal{G}\mathcal{I}_2^*]\hat{m}_\ell(\hat{r}_{K_1}-1)}{\hat{r}_{K_1}}-\frac{4Re[\mathcal{G}\mathcal{J}_2^*]\hat{m}_\ell(\hat{r}_{K_1}-1)}{\hat{r}_{K_1}}\Bigg\} \qquad (2.37)
\end{aligned}$$

The normalized differential forward-backward asymmetry of the $\overline{B} \to \overline{K}_1\ell^+\ell^-$ decay is defined by

$$\mathcal{A}_{FB}(\hat{s}) = \frac{\int_0^1 \Gamma(\hat{s},\cos(\theta))d\cos(\theta) - \int_{-1}^0 \Gamma(\hat{s},\cos(\theta))d\cos(\theta)}{\int_0^1 \Gamma(\hat{s},\cos(\theta))d\cos(\theta) + \int_{-1}^0 \Gamma(\hat{s},\cos(\theta))d\cos(\theta)} \qquad (2.38)$$

Using the definition mentioned above we calculate the normalized differential forward-backward asymmetry(A_{FB}). The result is as follows:

$$\begin{aligned}
\mathcal{A}_{FB}(\hat{s}) =& \frac{v\sqrt{\lambda}}{\hat{r}_{K_1}\Delta}\Bigg\{2(Re[\mathcal{A}\mathcal{F}^*]+Re[\mathcal{B}\mathcal{E}^*])\hat{r}_{K_1}\hat{s}+\hat{m}_\ell Re[\mathcal{B}(\mathcal{I}_1+\mathcal{J}_1)^*](-1+\hat{r}_{K_1}+\hat{s}) \\
&+\hat{m}_\ell Re[\mathcal{C}(\mathcal{I}_1+\mathcal{J}_1)^*]\lambda\Bigg\} \qquad (2.39)
\end{aligned}$$

Note that the pseudoscalar structure present in the decay amplitude(eq. (2.20)) can affect the branching ratio, the same structure doesn't contribute to the expression for the A_{FB}.

Parameter	Value
$\alpha_s(m_Z)$	0.119
α_{em}	1/129
$m_{K_1(1270)}$	1.270 (GeV) [42]
$m_{K_1(1400)}$	1.403 (GeV) [42]
$m_{K_{1A}}$	1.31 (GeV) [44]
$m_{K_{1B}}$	1.34 (GeV) [44]
m_b	4.8 (GeV)
m_μ	0.106 (GeV)
m_τ	1.780 (GeV)

Table 2. Input parameters.

F	$F(0)$	a	b	F	$F(0)$	a	b
$V_1^{BK_{1A}}$	0.34 ± 0.07	0.635	0.211	$V_1^{BK_{1B}}$	$-0.29^{+0.08}_{-0.05}$	0.729	0.074
$V_2^{BK_{1A}}$	0.41 ± 0.08	1.51	1.18	$V_2^{BK_{1B}}$	$-0.17^{+0.05}_{-0.03}$	0.919	0.855
$V_0^{BK_{1A}}$	0.22 ± 0.04	2.40	1.78	$V_0^{BK_{1B}}$	$-0.45^{+0.12}_{-0.08}$	1.34	0.690
$A^{BK_{1A}}$	0.45 ± 0.09	1.60	0.974	$A^{BK_{1B}}$	$-0.37^{+0.10}_{-0.06}$	1.72	0.912
$T_1^{BK_{1A}}$	$0.31^{+0.09}_{-0.05}$	2.01	1.50	$T_1^{BK_{1B}}$	$-0.25^{+0.06}_{-0.07}$	1.59	0.790
$T_2^{BK_{1A}}$	$0.31^{+0.09}_{-0.05}$	0.629	0.387	$T_2^{BK_{1B}}$	$-0.25^{+0.06}_{-0.07}$	0.378	-0.755
$T_3^{BK_{1A}}$	$0.28^{+0.08}_{-0.05}$	1.36	0.720	$T_3^{BK_{1B}}$	-0.11 ± 0.02	-1.61	10.2

Table 3. Formfactors for $B \to K_{1A}, K_{1B}$ transitions obtained in the LCQSR calculation [43] are fitted to the 3-parameter form in eq. (3.1).

Thus, the study of A_{FB} is complimentary to the study of branching ratio in order to extract the information about the nature of interactions in SUSY models.

3 Numerical results

In this section, we present the branching ratio and FB asymmetry for the $B \to K^*\ell^+\ell^-$ and the $B \to K_1(1270)\ell^+\ell^-$ decay for muon and tau leptons. The main input parameters are the form factors for which we use the results of light cone QCD sum rules(LCQCD) [43]. We use the parameters given in tables 2 and 3 in our numerical analysis. The values of the form factors at $q^2 = 0$ are given in table 3 [43]

The best fit for the q^2 dependence of the form factors can be written in the following form:

$$f_i(\hat{s}) = \frac{f_i(0)}{1 - a_i\hat{s} + b_i\hat{s}^2}\,, \tag{3.1}$$

The values of the parameters $f_i(0)$, a_i and b_i are given in table 3.

The mixing angle θ_{K_1} was estimated to be $|\theta_{K_1}| \approx 34^\circ \vee 57^\circ$ in ref. [45], $35^\circ \leq |\theta_{K_1}| \leq 55^\circ$ in ref. [46], $|\theta_{K_1}| = 37^\circ \vee 58^\circ$ in ref. [47], and $\theta_{K_1} = -(34\pm13)^\circ$ in [20, 48]. In this study, we use the results of ref. [20, 48] for numerical calculations, where we take $\theta_{K_1} = -34^\circ$.

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The new Wilson coefficients C_{Q_1} and C_{Q_2} describes in terms of masses of sparticles i.e., chargino-up-type squark and NHBs, $\tan(\beta)$ which is defined as the ratio of the two vacuum values of the 2 neutral Higgses and μ which has the dimension of a mass, corresponding to a mass term mixing the 2 Higgses doublets. It is obvious from eqs. (2.2) and (2.3) that C_{Q1} and C_{Q2} can reach a value of about one only when $\tan\beta$ is large enough($\beta \geq 20$) due to smallness of $m_b m_\ell / m_h^2 (h = h^0, A^0)$. Also, a magnitude of about one is consistent with the $B \to K^*\mu^+\mu^-$ rate and rate difference of electron channel to muon channel(R) at low dileptonic invariant mass region($1 \leq q^2 \leq 6\text{GeV}^2$) [49]. A similar result is found that in an MSSM-inspired scenario with large $\tan(\beta)$ and neutral Higgs exchange C_{Q1} and C_{Q2} are about one in magnitude which is consistent with current data [49, 50]. Note that μ can be positive or negative. Depending on the magnitude and sign of these parameters one can consider many options in the parameter space, but experimental results i.e., the rate of $b \to s\gamma$, $B \to K^*\mu^+\mu^-$ and $b \to s\ell^+\ell^-$ constrain us to consider the following options

- SUSY I: μ takes negative value, C_7 changes its sign and contribution of NHBs are neglected.
- SUSY II: $\tan(\beta)$ takes large values while the mass of superpartners are small i.e., few hundred GeV.
- SUSY III: $\tan(\beta)$ is large and the masses of superpartners are relatively large, i.e., about 450 GeV or more.

The numerical values of Wilson coefficients used in our analysis are modified the results of ref. [51, 52] according to the experimantal results obtained by BELLE collaboration [49] and those of ref. [50]. The numerical values of Wilson coefficients are collected in tables 4, and 5.

Moreover, in the absence of real experimental constraints on the FCNC modes into taus we could employ much larger Wilson coefficients (hence, SUSY effects) than we presented in tables 4, and 5, since the Yukawa-driven Higgs coupling implies that $C_Q^\tau = m_\tau / m_\mu C_Q^\mu$.

The numerical results for the decay rates and A_{FB}'s are presented in figures 1-2 and 4-5. Moreover, considering the new physics that couples to the mass of the lepton via the scalar and pseudoscalar type interactions clearly indicates that the decay rate for electron and muon channel can be different. We define a new observable $R(q^2)$ as follows:

$$R(q^2) = \frac{(d\Gamma/dq^2)(B \to K_1(1270)(K^*)e^+e^-)}{(d\Gamma/dq^2)(B \to K_1(1270)(K^*)\mu^+\mu^-)} \tag{3.2}$$

Figure 1 describes the differential decay rate of the $B \to K^*\mu^+\mu^-$ and the $B \to K_1(1270)\mu^+\mu^-$, from which one can see that the supersymmetric effects are quite significant (about twice of SM) for SUSY I and SUSY II models in the low momentum transfer regions, whereas these effects are small for SUSY III case. The reason for the increase of differential decay width in SUSY I model is the relative change in the sign of C_7^{eff} which gives dominant contribution in the low momentum transfer regions (look at the factor of $1/q^2$ in the eq. (2.1), while the large change in SUSY II model is owing to the contribution of the

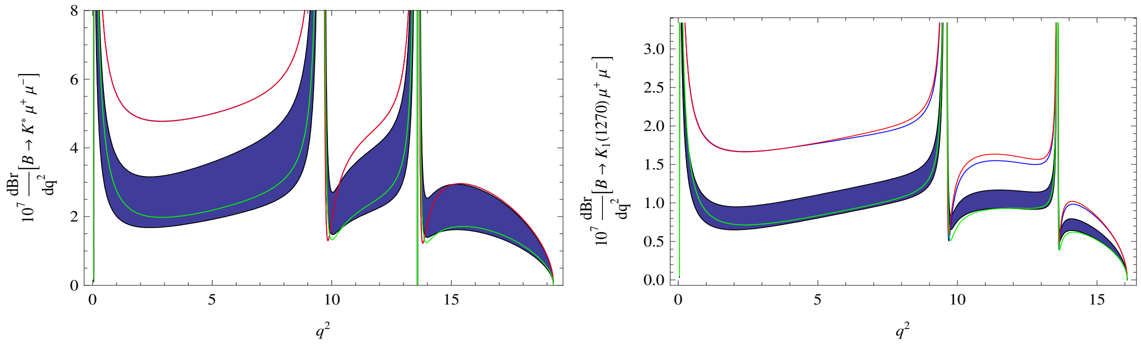

Fig.(1a) Fig.(1b)

Figure 1. Branching ratio of the $B \to K^*\mu^+\mu^-$ decay and the $B \to K_1(1270)\mu^+\mu^-$ decay. The black, blue, red and green lines correspond to SM, SUSY I, SUSY II, SUSY III models, respectively. The blue bound of the SM is created by the theoretical errors among the formfactors.

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NHBs. Furthermore, it can be seen that in the high q^2 region the SUSY effects are much more distinguishable for $K_1(1270)$ channel than K^* channel. The same effects can also be seen for the τ channel (see figure 4). Figure 2 describes the A_{FB} of the $B \to K^*\mu^+\mu^-$ and the $B \to K_1(1270)\mu^+\mu^-$, from which one can see that except SUSY III the supersymmetric effects are drastic in the low momentum transfer regions. In SUSY I and SUSY II models, the sign of $C_7^{\rm eff}$ and $C_9^{\rm eff}$ become the same, hence, the zero point of the A_{FB}'s disappears. Though, in the SUSY III model A_{FB} passes from the zero but this zero position shifts to the right from that of the SM value due to the contribution from the NHBs. A_{FB} is suppressed with the supersymmetric effects fot tau channel. The suppression is much more in the SUSY II model than the others (see figure 4). Figure 3 shows the dependency of R in terms of q^2 for various SUSY scenarios for $q^2 \geq 4m_\ell^2$ region. The study of rate difference of electron channel to muon channel as it can be seen is complimentary to the studies of other observables. While SUSY I and SUSY II are approximately coincide with each other in the study of branching ratio and A_{FB}, those models can be distinguished by studying the R (see figure 3). Furthermore, SUSY III lies in the theoretical error bounds of SM when looking at branching ratio and A_{FB} (see figures 1,2), but SM and SUSY III show different behavior in the R (see figure 3). The SUSY effects are larger for $K_1(1270)$ channel than K^* channel. The total or integrated branching ratio and the averaged forward-backward asymmetry for definite region are defined as

$$\mathcal{B}_r = \int \frac{d\mathcal{B}}{ds} ds,$$
$$\langle A_{FB} \rangle = \frac{\displaystyle\int A_{FB} \frac{d\mathcal{B}}{ds} ds}{\mathcal{B}_r}\,. \tag{3.3}$$

.

The integrated values of observables at low and high q^2 regions are calculated. The results of calculation and experimental values are depicted by tables 6-9. The results indicate that firstly, the chosen values of Wilson coefficients are consistent with measured rate of $B \to K^*\mu^+\mu^-$ at $1 \leq q^2 \leq 6\mathrm{GeV}^2$, secondly, the manifestation of the SUSY effects

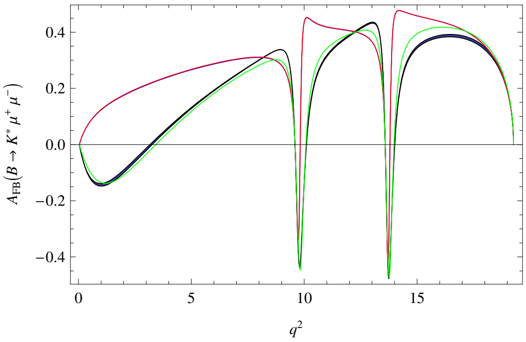

Figure (2a)

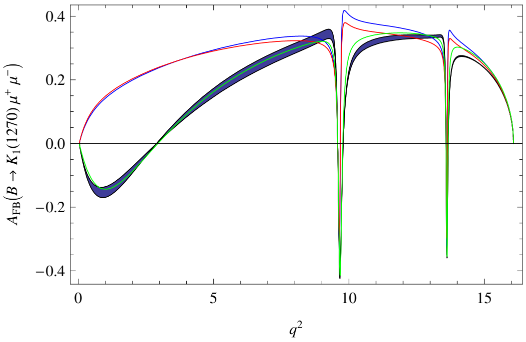

Figure (2b)

Figure 2. A_{FB} of the $B \to K^* \mu^+ \mu^-$ decay and the $B \to K_1(1270)\mu^+\mu^-$ decay. The black, blue, red and green lines correspond to the SM, SUSY I, SUSY II, SUSY III models, respectively. The blue bound of the SM is created by the theoretical errors among the formfactors.

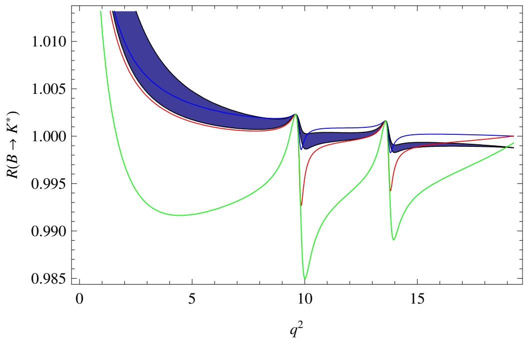

Figure (3a)

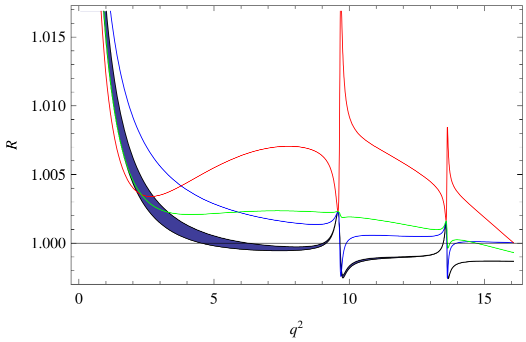

Figure (3b)

Figure 3. The rate difference of the electron channel to the muon channel for the $B \to K^*$ Figure (3a) and the $B \to K_1(1270)$ Figure (3b) transitions when $q^2 \geq 4m_\mu^2$ region. The blue bound of the SM is created by the theoretical errors among the formfactors.

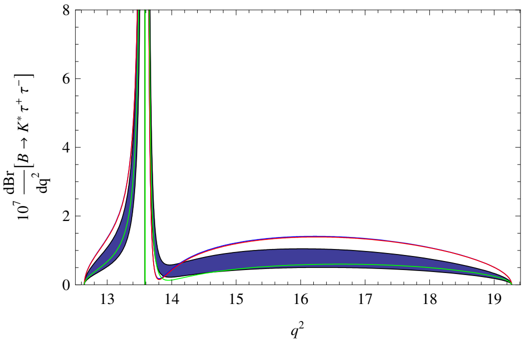

Figure (4a)

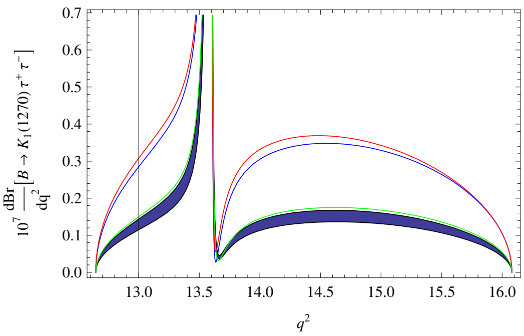

Figure (4b)

Figure 4. The same as figure 1 but for τ channel

Wilson Coefficients	$C_7^{\rm eff}$	C_9	C_{10}
SM	-0.313	4.334	-4.669
SUSY I	$+0.3756$	4.7674	-3.7354
SUSY II	$+0.3756$	4.7674	-3.7354
SUSY III	-0.3756	4.7674	-3.7354

Table 4. Wilson Coefficients in SM and different SUSY models without NHBs contributions.

Wilson Coefficients	C_{Q_1}	C_{Q_2}
SM	0	0
SUSY I	0	0
SUSY II	$0.5\,(16.5)$	$-0.5\,(-16.5)$
SUSY III	$1.2\,(4.5)$	$-1.2\,(-4.5)$

Table 5. Wilson coefficient corresponding to NHBs contributions within SUSY I, II and III models [51]. The values in the bracket are for the τ.

	$B \to K^*\mu^+\mu^-$	$B \to K_1\mu^+\mu^-$
$SM\ \mathcal{B}(10^{-7})$	$1.21^{+0.35}_{-0.39}$	$0.21^{+0.02}_{-0.02}$
SUSY I $\mathcal{B}(10^{-7})$	2.273	1.23
SUSY II $\mathcal{B}(10^{-7})$	2.270	1.2
SUSY III $\mathcal{B}(10^{-7})$	0.980	0.20
Exp. $\mathcal{B}(10^{-7})$	$1.49^{+0.45}_{-0.40} \pm 0.12$ [49]	–

Table 6. Experimentally measured values and integrated values of branching ratio at low dileptonic invariant mass region($1 \leq q^2 \leq 6\mathrm{GeV}^2$).

for $K_1(1270)$ channel are proper than K^* channel. To sum up, we study the semileptonic rare the $B \to K^*\ell^+\ell^-$ and the $B \to K_1(1270)\ell^+\ell^-$ decays in the MSSM with R parity. We show that the branching ratio and A_{FB} are very sensitive to the SUSY parameters. The branching ratio is enhanced up to one order of magnitude with respect to the corresponding SM values. The magnitude and sign of A_{FB} show quite a significant discrepancy with respect to the SM values. We also find that the study of rate difference of electron channel to muon channel R can be complimentary to the studies of branching ratio and A_{FB}. Also, it is found that the study of the $B \to K_1(1270)\ell^+\ell^-$ decay for SUSY effects can be proper than the $B \to K^*\ell^+\ell^-$ decay. Since, deviations for the $B \to K_1(1270)\ell^+\ell^-$ decay are greater than the $B \to K^*\ell^+\ell^-$ decay. The results of this study can be used to indirect search for the various SUSY effects in future planned experiments at LHC.

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	$B \to K^* \mu^+ \mu^-$	$B \to K_1 \mu^+ \mu^-$	$B \to K^* \tau^+ \tau^-$	$B \to K_1 \tau^+ \tau^-$
$SM\ \mathcal{B}(10^{-7})$	$0.158^{+0.004}_{-0.0004}$	$0.21^{+0.02}_{-0.02}$	$0.11^{+0.01}_{-0.01}$	$0.0185^{+0.015}_{-0.015}$
SUSY I $\mathcal{B}(10^{-7})$	0.181	1.23	0.083	0.0427
SUSY II $\mathcal{B}(10^{-7})$	0.184	1.2	0.086	0.0441
SUSY III $\mathcal{B}(10^{-7})$	0.173	0.20	0.12	0.0218

Table 7. Integrated values of branching ratio at high dileptonic invariant mass region($q^2 \geq 14.5\text{GeV}^2$).

	$B \to K^* \mu^+ \mu^-$	$B \to K_1 \mu^+ \mu^-$
SM	$0.00051^{+0.0001}_{-0.2}$	$0.0091^{+0.002}_{-0.002}$
SUSY I	0.0007	0.0184
SUSY II	0.0007	0.0183
SUSY III	0.00043	0.0077

Table 8. Averaged values of A_{FB} at low dileptonic invariant mass region($1 \leq q^2 \leq 6\text{GeV}^2$).

	$B \to K^* \mu^+ \mu^-$	$B \to K_1 \mu^+ \mu^-$	$B \to K^* \tau^+ \tau^-$	$B \to K_1 \tau^+ \tau^-$
SM	$0.326^{+0.0001}_{-0.0001}$	$0.195^{+0.005}_{-0.005}$	$0.236^{+0.0004}_{-0.0004}$	$0.125^{+0.0005}_{-0.0005}$
SUSY I	0.373	0.206	0.174	0.0817
SUSY II	0.374	0.205	0.181	0.0637
SUSY III	0.359	0.205	0.249	0.0985

Table 9. Averaged values of A_{FB} at high dileptonic invariant mass region($q^2 \geq 14.5\text{GeV}^2$).

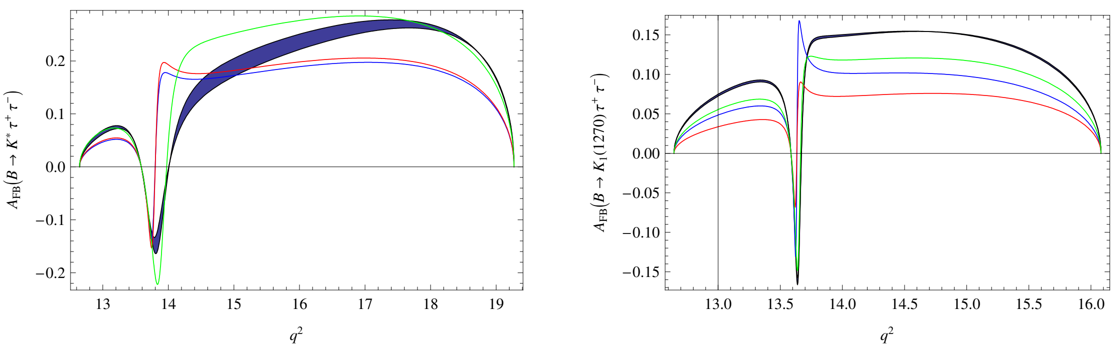

Figure (5a) Figure (5b)

Figure 5. The same as figure 2 but for τ channel

Acknowledgments

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