Forward-backward Asymmetry and Branching Ratio of $B \to K_1 \ell^+ \ell^-$ 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, $^3P_1(K_{1A}^1)$ and $^1P_1(K_{1B}^1)$. Taking into account this mixture, the forward-backward asymmetry and branching ratio of $B \to K_1(1270, 1400) \ell^+ \ell^-$ transitions are studied in the framework of different supersymmetric models. It is found that the results have considerable deviation from the standard model predictions. Any measurement of these physical observables and their comparison with the results obtained in this paper can give useful information about the nature of interactions beyond the standard model.

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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 $b \rightarrow s$ 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 $B \rightarrow K_1 \ell^+ \ell^-$ transition proceeds via the FCNC transition of $b \rightarrow s$ at quark level. $b \rightarrow s$ 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 $B \rightarrow X_s \ell^+ \ell^-$ decay is observed in BaBar [6] and Belle collaborations. These collaborations have also announced the measuring exclusive modes $B \rightarrow K \ell^+ \ell^-$ [7, 8, 9] and $B \rightarrow K^* \ell^+ \ell^-$ [10]. The obtained experimental results on these transitions are in a good consistency with theoretical predictions [11, 12, 13] the results of which can be used to constrain the new physics (NP) effects.

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 $B \rightarrow K_1 \ell^+ \ell^-$ 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[14]

$$
\begin{pmatrix}
\langle \bar{K}_1(1270) \rangle \\
\langle \bar{K}_1(1400) \rangle \\
\end{pmatrix} = M
\begin{pmatrix}
\langle \bar{K}_1^{A} \rangle \\
\langle \bar{K}_1^{B} \rangle \\
\end{pmatrix}, \quad \text{with} \quad M = \begin{pmatrix}
\sin \theta_{K_1} & \cos \theta_{K_1} \\
\cos \theta_{K_1} & -\sin \theta_{K_1} \\
\end{pmatrix}.
$$

The branching ratio of the $K_1(1400)$ case is smaller than the $K_1(1270)$ [14], so we consider only $B \rightarrow K_1(1270) \ell^+ \ell^-$. Note that lepton polarization and angular distribution of this decay in the frame work of SM has recently been studied in Refs. [15, 16].

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

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 model:

$$\mathcal{H}_{\text{eff}} = \frac{G_F \alpha m_b}{2\sqrt{2\pi}} \left[ C_9^{\text{eff}}(m_b) s \gamma_\mu(1 - \gamma_5) b \bar{\ell} \gamma_\mu \ell + C_{10}(m_b) s \gamma_\mu(1 - \gamma_5) b \bar{\ell} \gamma_\mu \gamma_5 \ell \right. $$

$$ \left. - 2m_b C_7(m_b) \frac{1}{q^2} s \sigma_{\mu\nu} q^\nu(1 + \gamma_5) b \bar{\ell} \gamma_\mu \ell + C_{Q_1} s(1 + \gamma_5) b \bar{\ell} \ell + C_{Q_2} s(1 + \gamma_5) b \bar{\ell} \gamma_5 \ell \right],$$

(2)

where $C_i$ are Wilson coefficients and the contributions of SUSY model are 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\cite{17, 18}. The $C_i$ 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\cite{19}–\cite{26}. $C_9^{\text{eff}}(s) = C_9 + Y(s)$, where $Y(s) = Y_{\text{pert}}(s) + Y_{\text{LD}}$ contains both the perturbative part $Y_{\text{pert}}(s)$ and long-distance part $Y_{\text{LD}}(s)$. $Y(s)_{\text{pert}}$ is given by\cite{19}

$$Y_{\text{pert}}(s) = g(m_c, s) c_0$$

$$- \frac{1}{2} g(1, s)(4c_3 + 4c_4 + 3c_5 + c_6) - \frac{1}{2} g(0, s)(c_3 + 3c_4)$$

$$+ \frac{2}{g}(3c_3 + c_4 + 3c_5 + c_6),$$

(3)

and the function $g(x, y)$ is defined in\cite{19}. Here, $c_1 - c_6$ are the Wilson coefficients in the leading logarithmic approximation. The relevant Wilson coefficients are given in Refs.\cite{27}. $Y(s)_{\text{LD}}$ involves $B \to K_1 V(\bar{c}c)$ resonances\cite{20}, where $V(\bar{c}c)$ are the vector charmonium states. We follow Refs.\cite{20, 28} and set

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

(5)

where $\hat{\Gamma}_{\text{tot}}^V \equiv \Gamma_{\text{tot}}^V / m_B$ and $\kappa_V$ takes different value for different exclusive semileptonic decay. The relevant properties of vector charmonium states are summarized in Table 1.

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), 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.

The $\mathcal{B}(p_B) \to \overline{K}_1(p_{K_1}, \lambda)$ form factors are defined by\cite{14}

$$\langle \overline{K}_1(p_{K_1}, \lambda) | \bar{s} \gamma_\mu(1 - \gamma_5) b | \mathcal{B}(p_B) \rangle$$
Table 1: Masses, total decay widths and branching fractions of dilepton decays of vector charmonium states [29].

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

\[
\begin{align*}
&= -i \frac{2}{m_B + m_{K_1}} \epsilon_{\mu\nu\rho\sigma}^{*,\nu}(p_B)^{\rho}(p_{K_1})^{\sigma} A_{1}^{K_1}(q^2) \\
&= \left[ (m_B + m_{K_1}) \epsilon_{\mu}^{(\lambda)*} V_{1}^{K_1}(q^2) - (p_B + p_{K_1})_{\mu} \epsilon_{(\lambda)}(p_B) \frac{V_{2}^{K_1}(q^2)}{m_B + m_{K_1}} \right] \\
&+ 2m_{K_1} \epsilon_{(\lambda)}^{*} \cdot p_B \frac{q_{\mu}}{q^2} V_{3}^{K_1}(q^2) - V_{0}^{K_1}(q^2), \tag{6}
\end{align*}
\]

where $q \equiv p_B - p_{K_1} = p_{e^+} + p_{e^-}$. By multiplying both sides of Eq. (6) with $q^\mu$, one can obtain the expression in terms of form factors for $\langle K_1(p_{K_1}, \lambda)|s(1 + \gamma_5)b|\overline{B}(p_B)\rangle$

\[
\begin{align*}
&= \frac{1}{m_b + m_s} \left\{ \\
&\quad - \left[ (m_B + m_{K_1}) \epsilon_{(\lambda)}^{*} q V_{1}^{K_1}(q^2) - (m_B - m_{K_1}) \epsilon_{(\lambda)} p_B V_{2}^{K_1}(q^2) \right] \\
&\quad + 2m_{K_1} \epsilon_{(\lambda)}^{*} \cdot p_B \left[ V_{3}^{K_1}(q^2) - V_{0}^{K_1}(q^2) \right] \right\}, \tag{8}
\end{align*}
\]

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 [14]):

\[
\begin{align*}
\left( \frac{\langle K_1(1270)|s(1 - \gamma_5)b|\overline{B}\rangle}{\langle K_1(1400)|s(1 - \gamma_5)b|\overline{B}\rangle} \right) &= M \left( \frac{\langle K_{1A}|s(1 - \gamma_5)b|\overline{B}\rangle}{\langle K_{1B}|s(1 - \gamma_5)b|\overline{B}\rangle} \right) ;
\left( \frac{\langle K_1(1270)|s(1 + \gamma_5)b|\overline{B}\rangle}{\langle K_1(1400)|s(1 + \gamma_5)b|\overline{B}\rangle} \right) &= M \left( \frac{\langle K_{1A}|s(1 + \gamma_5)b|\overline{B}\rangle}{\langle K_{1B}|s(1 + \gamma_5)b|\overline{B}\rangle} \right) ,
\end{align*}
\]
where the mixing matrix $M$ being given in Eq. (1) the formfactors $A^{K_1}, V^{K_1}_{0,1,2}$ and $T^{K_1}_{1,2,3}$ can be written as follows:

\[
\begin{align*}
\left( \frac{A^{K_1(1270)}}{(m_B + m_{K_1(1270)})} \right) &= M \left( \frac{A^{K_1A}}{(m_B + m_{K_1A})} \right), \\
\left( \frac{A^{K_1(1400)}}{(m_B + m_{K_1(1400)})} \right) &= M \left( \frac{A^{K_1B}}{(m_B + m_{K_1B})} \right), \\
\left( \frac{(m_B + m_{K_1(1270)})V_1^{K_1(1270)}}{(m_B + m_{K_1(1400)})V_1^{K_1(1400)}} \right) &= M \left( \frac{V_1^{K_1A}}{(m_B + m_{K_1A})} \right), \\
\left( \frac{(m_B + m_{K_1(1270)})V_0^{K_1(1270)}}{(m_B + m_{K_1(1400)})V_0^{K_1(1400)}} \right) &= M \left( \frac{V_0^{K_1A}}{(m_B + m_{K_1A})} \right), \\
\left( \frac{T_1^{K_1(1270)}}{T_1^{K_1(1400)}} \right) &= M \left( \frac{T_1^{K_1A}}{T_1^{K_1B}} \right), \\
\left( \frac{(m_B^2 - m_{K_1(1270)}^2)T_2^{K_1(1270)}}{(m_B^2 - m_{K_1(1400)}^2)T_2^{K_1(1400)}} \right) &= M \left( \frac{T_2^{K_1A}}{(m_B^2 - m_{K_1B}^2)T_2^{K_1B}} \right), \\
\left( \frac{T_3^{K_1(1270)}}{T_3^{K_1(1400)}} \right) &= M \left( \frac{T_3^{K_1A}}{T_3^{K_1B}} \right),
\end{align*}
\]

where it is supposed that $p^\mu_{K_1(1270),K_1(1400)} \simeq p^\mu_{K_1A} \simeq p^\mu_{K_1B}[14]$. These formfactors within light-cone QCD sum rule (LCQSR) are estimated in [30].

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

\[
\mathcal{M} = \frac{G_{F\alpha_em}}{2\sqrt{2\pi}} V_{ts}^* V_{tb} m_B \cdot (-i) \left\{ T^{(K_1)1}_\mu \bar{\ell}_\gamma \mu \ell + T^{(K_1)2}_\mu \bar{\ell}_\gamma \mu \gamma_5 \ell + T^{(K_1)3} \bar{\ell} \ell + T^{(K_1)4} \bar{\ell} \gamma_5 \ell \right\},
\]

where

\[
\begin{align*}
T^{(K_1)1}_\mu &= A^{K_1}(\hat{s}) \epsilon_{\mu
u\rho\sigma} \epsilon^{*\nu\rho\sigma} \hat{p} \hat{q} - iB^{K_1}(\hat{s}) \epsilon_\mu, \\
T^{(K_1)2}_\mu &= \mathcal{E}^{K_1}(\hat{s}) \epsilon_{\mu
u\rho\sigma} \epsilon^{*\nu\rho\sigma} \hat{p} \hat{q} - i\mathcal{F}^{K_1}(\hat{s}) \epsilon_\mu, \\
T^{(K_1)3} &= i\mathcal{I}^{K_1}_1(\hat{s}) \frac{\epsilon^{(\lambda)*}}{1 + \hat{m}_s} + i\mathcal{J}^{K_1}_1(\hat{s}) \frac{\epsilon^{(\lambda)*}}{1 + \hat{m}_s}, \\
T^{(K_1)4} &= i\mathcal{I}^{K_1}_2(\hat{s}) \frac{\epsilon^{(\lambda)*}}{1 + \hat{m}_s} + i\mathcal{J}^{K_1}_2(\hat{s}) \frac{\epsilon^{(\lambda)*}}{1 + \hat{m}_s},
\end{align*}
\]

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

\[
A^{K_1}(\hat{s}) = \frac{2}{1 + \sqrt{r_{K_1}}} \epsilon^{eff}_{K_1}(\hat{s}) A^{K_1}(\hat{s}) + \frac{4\hat{m}_b}{\hat{s}} \epsilon^{eff}_{\gamma T} T^{K_1}_1(\hat{s}),
\]

(23)
with \( \hat{v} \) given by

\[
D_{K_i}(s) = \frac{1}{1 - \hat{r}_{K_i}} \left[ (1 - \sqrt{\hat{r}_{K_i}}) c_{9}^{\text{eff}}(\hat{s}) V_{1}^{K_1}(\hat{s}) + 2\hat{m}_b c_{7}^{\text{eff}} \left( T_{3}^{K_1}(\hat{s}) + \frac{1 - \sqrt{\hat{r}_{K_i}}}{\hat{s}} T_{2}^{K_1}(\hat{s}) \right) \right],
\]

(24)

\[
C_{K_i}(s) = \frac{1}{1 - \hat{r}_{K_i}} \left[ (1 - \sqrt{\hat{r}_{K_i}}) c_{9}^{\text{eff}}(\hat{s}) V_{2}^{K_1}(\hat{s}) + \hat{2} \hat{m}_b c_{7}^{\text{eff}} \left( T_{3}^{K_1}(\hat{s}) + \frac{1 - \sqrt{\hat{r}_{K_i}}}{s^2} T_{2}^{K_1}(\hat{s}) \right) \right],
\]

(25)

\[
D_{K_1}(s) = \frac{1}{\hat{s}} c_{9}^{\text{eff}}(\hat{s}) \left\{ (1 + \sqrt{\hat{r}_{K_i}}) V_{1}^{K_1}(\hat{s}) - (1 - \sqrt{\hat{r}_{K_i}}) V_{2}^{K_1}(\hat{s}) - 2\sqrt{\hat{r}_{K_i}} V_{0}^{K_1}(\hat{s}) \right\}
\]

(26)

\[
E_{K_i}(s) = \frac{2}{1 + \sqrt{\hat{r}_{K_i}}} c_{10} A_{K_i}(\hat{s}),
\]

(27)

\[
F_{K_i}(s) = (1 + \sqrt{\hat{r}_{K_i}}) V_{1}^{K_1}(\hat{s}),
\]

(28)

\[
G_{K_i}(s) = \frac{1}{1 + \sqrt{\hat{r}_{K_i}}} c_{10} V_{2}^{K_1}(\hat{s}),
\]

(29)

\[
H_{K_i}(s) = \frac{1}{\hat{s}} c_{10} \left[ (1 + \sqrt{\hat{r}_{K_i}}) V_{1}^{K_1}(\hat{s}) - (1 - \sqrt{\hat{r}_{K_i}}) V_{2}^{K_1}(\hat{s}) - 2\sqrt{\hat{r}_{K_i}} V_{0}^{K_1}(\hat{s}) \right],
\]

(30)

\[
I_{K_1}(s) = -C_{Q_1}(1 + \sqrt{\hat{r}_{K_i}}) V_{1}^{K_1}(\hat{s})
\]

(31)

\[
J_{K_1}(s) = C_{Q_1} \{ (1 + \sqrt{\hat{r}_{K_i}}) V_{2}^{K_1}(\hat{s}) + 2\sqrt{\hat{r}_{K_i}} [V_{3}^{K_1}(\hat{s}) - V_{0}^{K_1}(\hat{s})] \}
\]

(32)

\[
K_{2}(s) = T_{2}^{K_1}(\hat{s})(C_{Q_2} \rightarrow C_{Q_1}), \quad J_{2}^{K_1}(s) = J_{1}^{K_1}(s)(C_{Q_2} \rightarrow C_{Q_1})
\]

(33)

with \( \hat{r}_{K_i} = m_{K_i}^2/m_{B}^2 \) and \( \hat{s} = q^2/m_{B}^2 \).

The dilepton invariant mass spectrum of the lepton pair for the \( B \rightarrow K_1 \ell^+ \ell^- \) decay is given by

\[
\frac{d \Gamma(B \rightarrow K_1 \ell^+ \ell^-)}{d \hat{s}} = \frac{G_{F}^{2} \alpha_{em}^{2} m_{B}^{5}}{212 \pi^{5}} |V_{tb} V_{ts}^{*}| ^{2} \nu \sqrt{\lambda} \Delta(\hat{s})
\]

(34)

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

\[
\Delta(\hat{s}) = \frac{8Re[\mathcal{F}_{K_i}^{*}] \hat{m}_{\ell}^{2} \lambda}{\hat{r}_{K_i}} + \frac{8Re[\mathcal{G}_{K_i}^{*}] \hat{m}_{\ell}^{2} (-1 + \hat{r}_{K_i}) \lambda}{\hat{r}_{K_i}} - \frac{8|\mathcal{H}|^{2} \hat{m}_{\ell}^{2} \hat{s} \lambda}{3 \hat{r}_{K_i}}
\]

\[
- \frac{2Re[\mathcal{B}^{*}](1 \hat{r}_{K_i} + \hat{s})(3 + 2\hat{r}_{K_i} - 6 \hat{s} + 3 \hat{s}^{2} - 6 \hat{r}_{K_i}(1 + \hat{s}) - v^{2} \lambda)}{3 \hat{r}_{K_i}}
\]

\[
- \frac{|C|^{2} \lambda (3 + 2\hat{r}_{K_i} - 6 \hat{s} + 3 \hat{s}^{2} - 6 \hat{r}_{K_i}(1 + \hat{s}) - v^{2} \lambda)}{3 \hat{r}_{K_i}}
\]

\[
- \frac{|G|^{2} \lambda (3 + 2\hat{r}_{K_i} + 12 \hat{m}_{\ell}^{2}(2 + 2\hat{r}_{K_i} - \hat{s}) - 6 \hat{s} + 3 \hat{s}^{2} - 6 \hat{r}_{K_i}(1 + \hat{s}) - v^{2} \lambda)}{3 \hat{r}_{K_i}}
\]

\[
+ \frac{|\mathcal{F}|^{2} (3 - 2\hat{r}_{K_i} + 6 \hat{r}_{K_i}(1 + 16 \hat{m}_{\ell}^{2} - 3 \hat{s}) + 6 \hat{s} - 3 \hat{s}^{2} + v^{2} \lambda)}{3 \hat{r}_{K_i}}
\]

\[
+ \frac{|\mathcal{B}|^{2} (3 - 2\hat{r}_{K_i} + 6 \hat{s} - 3 \hat{s}^{2} - 6 \hat{r}_{K_i}(-1 + 8 \hat{m}_{\ell}^{2} + 3 \hat{s}) + v^{2} \lambda)}{3 \hat{r}_{K_i}}
\]
In Tables 2 and 3 in our numerical analysis. The values of the form factors at decay for muon and tau leptons. The main input parameters are the form factors for which

In this section, we present the branching ratio and FB asymmetry for the

Note that the pseudoscalar structure existing in the decay amplitude(Eq. 18) can affect the

The normalized differential forward-backward asymmetry of the $B \to K_1 \ell^+ \ell^-$ decay is defined by

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

Note that the pseudoscalar structure existing in the decay amplitude(Eq. 18) can affect the branching ratio, the same structure don’t contribute to the expression of the FBA. Thus, the study of FBA 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 $B \to K_1(1270) \ell^+ \ell^-$ 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) [30]. 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[30]

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

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

(38)
Table 2: Input parameters

| Parameter | Value  |
|-----------|--------|
| $\alpha_s(m_Z)$ | 0.119  |
| $\alpha_{em}$ | 1/129  |
| $m_{K_1}(1270)$ | 1.270 (GeV) [29] |
| $m_{K_1}(1400)$ | 1.403 (GeV) [29] |
| $m_{K_1A}$ | 1.31 (GeV) [31] |
| $m_{K_1B}$ | 1.34 (GeV) [31] |
| $m_b$ | 4.8 (GeV) |
| $m_{\mu}$ | 0.106 (GeV) |
| $m_\tau$ | 1.780 (GeV) |

Table 3: Formfactors for $B \to K_{1A}, K_{1B}$ transitions obtained in the LCQSR calculation [30] are fitted to the 3-parameter form in Eq. (38).

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

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

The mixing angle $\theta_{K_1}$ was estimated to be $|\theta_{K_1}| \approx 34^\circ \sqrt{57^\circ}$ in Ref. [32], $35^\circ \leq |\theta_{K_1}| \leq 55^\circ$ in Ref. [33], $|\theta_{K_1}| = 37^\circ \sqrt{58^\circ}$ in Ref. [34], and $\theta_{K_1} = -(34 \pm 13)^\circ$ in [14, 35]. In this study, we use the results of Ref.[14, 35] for numerical calculations, where we take $\theta_{K_1} = -34^\circ$.

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 $\mu$ which has the dimension of a mass, corresponding to a mass term mixing the 2 Higgs doublets. Note that $\mu$ 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 $b \to s\gamma$ and $b \to s\ell^+\ell^-$ constrain us to consider the following options

- **SUSY I**: $\mu$ 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 borrowed from Ref. [36, 37] and collected in Tables 4, and 5.

The numerical results for the decay rates and FBAs are presented in Figs. 1-4. Fig. 1 describes the differential decay rate of $B \rightarrow K_1(1270)\mu^+\mu^-$, 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), while the large change in SUSY II model is owing to the contribution of the NHBs. The same effects can also be seen for the $\tau$ channel (see fig. 2). Fig. 3 describes the FBA of $B \rightarrow K_1(1270)\mu^+\mu^-$, 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 $C_7^{eff}$ and $C_9^{eff}$ become the same, hence, the zero point of the FBAs disappears. Though, in the SUSY III model FBA 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. FBA is suppressed with the supersymmetric effects. The suppression is much more in the SUSY II model than the others (see fig. 4).

| Wilson Coefficients | $C_7^{eff}$ | $C_9$ | $C_{10}$ |
|---------------------|-------------|-------|---------|
| 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             | 6.5 (16.5) | -6.5 (-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 [36]. The values in the bracket are for the $\tau$.

To sum up, we study the semileptonic rare $B \rightarrow K^*_1(1270)\ell^+\ell^-$ decay in the supersymmetric theories. We show that the branching ratio and FBA 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 FBA show quite a significant discrepancy with respect to the SM values. The results of this study can be used to indirect search for the SUSY effects in future planned experiments at LHC.
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Figure 1: Branching ratio of $B \rightarrow K^*_1(1270)\mu^+\mu^-$ decay. The black, blue, red and green lines correspond to SM, SUSY I, SUSY II, SUSY III models, respectively.

Figure 2: The same as Fig. 1 but for $\tau$ channel
Figure 3: FBA of $B \rightarrow K^{+}(1270)\mu^{+}\mu^{-}$ decay. The black, blue, red and green lines correspond to SM, SUSY I, SUSY II, SUSY III models, respectively.

Figure 4: The same as Fig. 3 but for $\tau$ channel