PHYSICAL REVIEW D **79,** 074007 (2009)

$B \to K_0^*(1430) l^+ l^-$ decays in supersymmetric theories

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(Received 9 February 2009; published 9 April 2009)

The weak decays of $B \to K_0^*(1430) l^+ l^-$ ($l = \mu, \tau$) are investigated in minimal supersymmetric standard model (MSSM) and also in supersymmetric (SUSY) SO(10) grand unified theory (GUT) models. Neutral Higgs bosons are the point of main focus in MSSM because they make quite a large contribution in exclusive $B \to X_s l^+ l^-$ decays at large $\tan\beta$ regions of parameter space of SUSY models, as part of SUSY contributions is proportional to $\tan^3\beta$. The analysis of decay rate, forward-backward asymmetries and lepton polarization asymmetries in $B \to K_0^*(1430) l^+ l^-$ show that the values of these physical observables are greatly modified by the effects of neutral Higgs bosons. In SUSY SO(10) grand unified model, the new physics contribution comes from the operators which are induced by the neutral Higgs boson penguins and also from the operators with chirality opposite to that of the corresponding standard model operators. SUSY SO(10) effects show up only in the decay $B \to K_0^* \mu^+ \mu^-$ where the transverse lepton polarization asymmetries deviate significantly from the SM value while the effects in the decay rate, forward-backward asymmetries, the longitudinal and normal lepton polarization asymmetries are very mild. The transverse lepton polarization asymmetry is almost zero in SM and in MSSM model, whereas it can reach to -0.3 in SUSY SO(10) GUT model which could be seen at the future colliders; hence this asymmetry observable can be used to discriminate between different SUSY models.

DOI: 10.1103/PhysRevD.79.074007 PACS numbers: 13.30.Ce, 11.30.Pb, 14.40.Nd

I. INTRODUCTION

It is generally believed that the standard model (SM) is one of the most successful theory of the second half of the last century as it has passed all the experimental tests carried out for its verifications. The only missing thing that is yet to be verified is the Higgs boson mass, which we hope will be measured at the CERN large hadron collider (LHC) in next couple of years. To test the SM indirectly, rare decays induced by flavor changing neutral currents (FCNCs) $b \to s l^+ l^-$ have become the main focus of the studies since the CLEO measurement of the radiative decay $b \to s\gamma$ [1]. In the SM these decays are forbidden at tree level and can only be induced via loop diagrams. Hence, such decays will provide useful information about the parameters of Cabbibo-Kobayashi-Maskawa (CKM) matrix [2,3] elements as well as the various hadronic form factors. In literature, there have been intensive studies on the exclusive decays $B \to P(V, A) l^+ l^-$ [4–10] both in the SM and beyond, where the notions P, V, and A denote the pseudoscalar, vector and axial vector mesons, respectively.

Despite all the success of the SM no one can say that it is the ultimate theory of nature as it has many open questions, such as the gauge hierarchy problem, origin of masses and Yukawa couplings, etc. It is known that supersymmetry (SUSY) is not only one of the strongest competitor of the SM but is also the most promising candidate of new physics. One direct way to search for SUSY is to discover SUSY particles at high energy colliders, but unfortunately, so far no SUSY particles have been found. Another way is to search for its effects through indirect methods. The measurement of invariant mass spectrum, forward-backward asymmetry and polarization asymmetries are the suitable tools to probe new physics effects. For most of the SUSY models, the SUSY contributions to an observable appear at loop level due to the R-parity conservation. Therefore, it has been realized for a long time that rare processes can be used as a good probe for the searches of SUSY, since in these processes the contributions of SUSY and SM arises at the same order in perturbation theory [11].

In other SUSY models, neutral Higgs bosons (NHBs) could contribute largely to the inclusive processes $B \to X_s l^+ l^-$, because part of the SUSY contributions is proportional to the $\tan^3\beta$ [12]. Subsequently, the physical observables, like branching ratio and forward-backward asymmetry, in the large $\tan\beta$ region of parameter space in SUSY models can be quite different from that in the SM. Motivated by the fact, similar effects in exclusive $B \to K(K^*) l^+ l^-$ decay modes are also investigated [11], where the analysis of decay rates, forward-backward asymmetries and polarization asymmetries of final state lepton indicates the significant role of NHBs. It is believed that the problem of neutrino oscillations can not be explained in the SM. To this purpose, the SUSY SO(10) grand unified theory (GUT) models [13] has been proposed in literature. In this model, there is a complex flavor nondiagonal down-type squark mass matrix element of 2nd and 3rd generations which is of the order one at the GUT scale. This can induce large flavor off-diagonal coupling such as the coupling of gluino to the quark and squark which belong to different generations. In general these couplings are complex and may contribute to the process of FCNCs. The above analysis of physical observables in $B \to K(K^*)l^+l^-$ decay is extended in SUSY SO(10) GUT model where it has been shown that the forward-backward asymmetries as well as the longitudinal and transverse decay widths of the said decays are sensitive to these NHBs effect in SUSY SO (10) GUT model which can be detected in the future B factories [14].

Along this line, we will investigate the exclusive decay $\bar{B}_0 \to K_0^*(1430)l\bar{l}$ ($l = \mu, \tau$), where $K_0^*(1430)$ is a scalar meson, both in the minimal supersymmetric standard model (MSSM) as well as in the SUSY SO(10) GUT model [13]. It is expected that we might doubly benefit from the study of such rare B decays, gain better understanding on low energy QCD and search for SUSY. The reason is obvious that in the dominant decay modes, the SM contribution is overwhelming and the SUSY effects, if exists, would be drawn in the background, however, for rare decay modes, situation would be different, where the SM contribution might be less significant, so that the SUSY effect would emerge to be distinguished from the SM background. On the QCD side, the priority to scrutinize the B decays involving the scalar meson, such as $K_0^*(1430)$, can be attributed to the fact that those mesons have the same quantum numbers as the vacuum and therefore can condensate into the vacuum as well as break a symmetry like a global chiral $U(N_f) \times U(N_f)$. Also, the nature of the scalar meson continues to be an intriguing problem at present, and the weak productions of them in the heavy hadron decays would allow us to probe the inner structure of these novel mesons in a even broader ground. Now, the main motivation to extend the SUSY analysis of $B \to K(K^*)l\bar{l}$ to $\bar{B}_0 \to K_0^*(1430)l\bar{l}$ ($l = \mu, \tau$) lies in the fact that because of the scalar meson in the final state, the dependence of various asymmetries on Wilson coefficients is different compared to the pseudoscalar and vector case. Hence it is possible that experimental results on $K_0^*(1430)$ can give us some insights that are not accessible in the studies of K or K^* modes.

In this paper, we evaluate the branching ratios, forward-backward asymmetries, lepton polarization asymmetries with special emphasis on the effects of NHBs in MSSM. It is known that different source of the vector current could manifest themselves in different regions of phase space. For low value of momentum transfer, the photonic penguin dominates, while the Z penguin and W box become important towards high value of momentum transfer [11]. In order to search the region of momentum transfer with large contributions from NHBs, the above decay in certain large $\tan\beta$ region of parameter space has been analyzed in SuperGravity (SUGRA) and M-theory inspired models [15]. We extend this analysis to the SUSY SO(10) GUT model [11], where there are some primed counterparts of the usual SM operators. For instance, the counterparts of usual operators in $B \to X_s\gamma$ decay are suppressed by m_s/m_b and consequently negligible in the SM because they have opposite chiralities. These operators are also suppressed in minimal flavor violating (MFV) models [16,17], however, in SUSY SO(10) GUT model their effects can be significant. The reason is that the flavor non-diagonal squark mass matrix elements are the free parameters, some of which have significant effects in rare decays of B mesons [18].

The main job of investigating the semileptonic B meson decay is to properly evaluate the hadronic matrix elements for $B \to K_0^*(1430)$, namely, the transition form factors, which are governed by the nonperturbative QCD dynamics. Several methods exist in the literature to deal with this problem, among which the QCD sum rules approach (QCDSR) [19,20] is a fully relativistic approach and well rooted in quantum field theory. However, short-distance expansion fails in nonperturbative condensate when applying the three-point sum rules to the computations of form factors in the large momentum transfer or large mass limit of heavy meson decays. As a marriage of standard QCDSR technique and theory of hard exclusive process, the light cone QCD sum rules (LCSR) [21–23] cure the problem of QCDSR applying to the large momentum transfer by performing the operator product expansion (OPE) in terms of twist of the relevant operators rather than their dimension [24]. Therefore, the principal discrepancy between QCDSR and LCSR consists in that nonperturbative vacuum condensates representing the long-distance quark and gluon interactions in the short-distance expansion are substituted by the light cone distribution amplitudes (LCDAs) describing the distribution of longitudinal momentum carried by the valence quarks of hadronic bound system in the expansion of transverse-distance between partons in the infinite momentum frame. An important advantage of LCSR is that it allows a systematic inclusion of both hard scattering effects and the soft contributions. Phenomenologically, LCSR has been widely applied to investigate the semileptonic decays of heavy hadrons [25–28], radiative hadronic decays [29–31] and nonleptonic two body decays of B meson [32–35].

In our numerical analysis for $\bar{B}_0 \to K_0^*(1430)$ decays, we shall use the results of the form factors calculated by LCSR approach in Ref. [36]. The values of the relevant Wilson coefficient for MSSM and SUSY SO(10) GUT models are borrowed from Ref. [11,14]. The effects of SUSY contributions to the decay rate and lepton polarization are also explored in this work. Our results show that the decay rates are quite sensitive to the NHBs contribution. The forward-backward asymmetry is zero in the SM for these decays because of the absence of the scalar-type coupling, therefore any nonzero value of the forward-backward asymmetry will give us indication of the new physics. Contrary to $B \to K\tau^+\tau^-$ decay the value of the forward-backward asymmetry for $\bar{B}_0 \to K_0^*(1430)\tau^+\tau^-$ is positive and is significant when the contributions of NHB's is large. Therefore the measurement of this observable will give us some clear indications of SUSY in these decays.

It is known that the hadronic uncertainties associated with the form factors and other input parameters have negligible effects on the lepton polarization asymmetries, therefore we have also studied these asymmetries in the SUSY models mentioned above and found that the effects of NHBs are quite significant in some regions of parameter space of SUSY. In $B \to K(K^*) l \bar{l}$ decays the normal lepton polarization is proportional to the imaginary part of the Wilson coefficients involved and is expected to be very small [11]. However, for $\bar{B}_0 \to K_0^*(1430) l \bar{l}$ it is proportional to the real part, hence the value is expected to be large in comparison to the previous case. Therefore, the experimental investigation of normal lepton polarization will give us some insights that are not accessible to the studies of $K(K^*)$ modes.

The paper is organized as follows. In Sec. II, we present the effective Hamiltonian for the semileptonic decay $B \to K_0^* l^+ l^-$. Section III contains the parametrizations and numbers of the form factors for the said decay using the LCSR approach. In Sec. IV we present the basic formulas of physical observables like decay rates, forward-backward asymmetries and polarization asymmetries of lepton in the above mentioned decay. Section V is devoted to the numerical analysis of these observables and the brief summary and concluding remarks are given in Sec. VI.

II. EFFECTIVE HAMILTONIAN

By integrating out the heavy degrees of freedom in the full theory, the general effective Hamiltonian for $b \to s l^+ l^-$ in SUSY SO(10) GUT model, can be written as [14]

$$H_{\rm eff} = -\frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^* \Bigg[\sum_{i=1}^{2} C_i(\mu) O_i(\mu) + \sum_{i=3}^{10} (C_i(\mu) O_i(\mu) + C_i'(\mu) O_i'(\mu)) + \sum_{i=1}^{8} (C_{Q_i}(\mu) Q_i(\mu) + C_{Q_i}'(\mu) Q_i'(\mu)) \Bigg], \quad (1)$$

where $O_i(\mu)$ ($i = 1, \ldots, 10$) are the four-quark operators and $C_i(\mu)$ are the corresponding Wilson coefficients at the energy scale μ [37]. Using renormalization group equations to resum the QCD corrections, Wilson coefficients are evaluated at the energy scale $\mu = m_b$. The theoretical uncertainties associated with the renormalization scale can be substantially reduced when the next-to-leading-logarithm corrections are included. The new operators $Q_i(\mu)$ ($i = 1, \ldots, 8$) come from the NHBs exchange diagrams, whose manifest forms and corresponding Wilson coefficients can be found in [38,39]. The primed operators are the counterparts of the unprimed operators, which can be obtained by flipping the chiralities in the corresponding unprimed operators. It is believed that the effects of the counterparts of usual chromomagnetic and electromagnetic dipole moment operators as well as semileptonic operators with opposite chirality are suppressed by m_s/m_b in the SM, but in SUSY SO(10) GUTs their effect can be significant, since δ_{23}^{dRR} can be as large as 0.5 [13,14]. Apart from this, δ_{23}^{dRR} can induce new operators as the counterparts of usual scalar operators in SUSY models due to NHB penguins with gluino-down type squark propagator in the loop. It is worth mentioning that these primed operators will appear only in SUSY SO(10) GUT model and are absent in SM and MSSM [11].

The explicit expressions of the operators responsible for $B \to K_0^*(1430) l^+ l^-$ transition are given by

$$\begin{aligned}
O_7 &= \frac{e^2}{16\pi^2} m_b (\bar{s} \sigma_{\mu\nu} P_R b) F^{\mu\nu}, \\
O_7' &= \frac{e^2}{16\pi^2} m_b (\bar{s} \sigma_{\mu\nu} P_L b) F^{\mu\nu} \\
O_9 &= \frac{e^2}{16\pi^2} (\bar{s} \gamma_\mu P_L b)(\bar{l} \gamma^\mu l), \\
O_9' &= \frac{e^2}{16\pi^2} (\bar{s} \gamma_\mu P_R b)(\bar{l} \gamma^\mu l) \\
O_{10} &= \frac{e^2}{16\pi^2} (\bar{s} \gamma_\mu P_L b)(\bar{l} \gamma^\mu \gamma_5 l), \\
O_{10}' &= \frac{e^2}{16\pi^2} (\bar{s} \gamma_\mu P_R b)(\bar{l} \gamma^\mu \gamma_5 l) \\
Q_1 &= \frac{e^2}{16\pi^2} (\bar{s} P_R b)(\bar{l} l), \\
Q_1' &= \frac{e^2}{16\pi^2} (\bar{s} P_L b)(\bar{l} l) \\
Q_2 &= \frac{e^2}{16\pi^2} (\bar{s} P_R b)(\bar{l} \gamma_5 l), \\
Q_2' &= \frac{e^2}{16\pi^2} (\bar{s} P_L b)(\bar{l} \gamma_5 l)
\end{aligned} \quad (2)$$

with $P_{L,R} = (1 \pm \gamma_5)/2$. In terms of the above Hamiltonian, the free quark decay amplitude for $b \to s l^+ l^-$ can be derived as [12]:

$$\begin{aligned}
\mathcal{M}(b \to s l^+ l^-) = -\frac{G_F \alpha}{\sqrt{2}\pi} V_{tb} V_{ts}^* \Big\{ & C_9^{\rm eff} (\bar{s} \gamma_\mu P_L b)(\bar{l} \gamma^\mu l) \\
& + C_{10} (\bar{s} \gamma_\mu P_L b)(\bar{l} \gamma^\mu \gamma_5 l) \\
& - 2 m_b C_7^{\rm eff} \Big(\bar{s} i \sigma_{\mu\nu} \frac{q^\nu}{s} P_R b \Big)(\bar{l} \gamma^\mu l) \\
& + C_{Q_1} (\bar{s} P_R b)(\bar{l} l) + C_{Q_2} (\bar{s} P_R b)(\bar{l} \gamma_5 l) \\
& + (C_i(m_b) \leftrightarrow C_i'(m_b)) \Big\},
\end{aligned} \quad (3)$$

where $s = q^2$ and q is the momentum transfer. The operator O_{10} can not be induced by the insertion of four-quark operators because of the absence of the Z boson in the effective theory. Therefore, the Wilson coefficient C_{10} does not renormalize under QCD corrections and hence it is independent on the energy scale. In addition to this, the above quark level decay amplitude can receive contributions from the matrix element of four-quark operators, $\sum_{i=1}^{6}\langle l^+l^-s|O_i|b\rangle$, which are usually absorbed into the effective Wilson coefficient $C_9^{\rm eff}(\mu)$, that one can decompose into the following three parts [40–46]

$$C_9^{\rm eff}(\mu) = C_9(\mu) + Y_{\rm SD}(z,s') + Y_{\rm LD}(z,s'),$$

where the parameters z and s' are defined as $z = m_c/m_b$, $s' = q^2/m_b^2$. $Y_{\rm SD}(z,s')$ describes the short-distance contributions from four-quark operators far away from the $c\bar{c}$ resonance regions, which can be calculated reliably in the perturbative theory. The long-distance contributions $Y_{\rm LD}(z,s')$ from four-quark operators near the $c\bar{c}$ resonance cannot be calculated from first principles of QCD and are usually parametrized in the form of a phenomenological Breit-Wigner formula making use of the vacuum saturation approximation and quark-hadron duality. We will neglect the long-distance contributions in this work because of the absence of experimental data on $B \to J/\psi K_0^*(1430)$. The manifest expressions for $Y_{\rm SD}(z,s')$ can be written as

$$\begin{aligned} Y_{\rm SD}(z,s') = {}& h(z,s')(3C_1(\mu) + C_2(\mu) + 3C_3(\mu) + C_4(\mu) \\ &+ 3C_5(\mu) + C_6(\mu)) - \frac{1}{2}h(1,s')(4C_3(\mu) \\ &+ 4C_4(\mu) + 3C_5(\mu) + C_6(\mu)) \\ &- \frac{1}{2}h(0,s')(C_3(\mu) + 3C_4(\mu)) + \frac{2}{9}(3C_3(\mu) \\ &+ C_4(\mu) + 3C_5(\mu) + C_6(\mu)), \end{aligned} \tag{4}$$

with

$$\begin{aligned} h(z,s') = {}& -\frac{8}{9}\ln z + \frac{8}{27} + \frac{4}{9}x - \frac{2}{9}(2+x)|1-x|^{1/2} \\ &\times \begin{cases} \ln\left|\frac{\sqrt{1-x}+1}{\sqrt{1-x}-1}\right| - i\pi & \text{for } x \equiv 4z^2/s' < 1 \\ 2\arctan\frac{1}{\sqrt{x-1}} & \text{for } x \equiv 4z^2/s' > 1 \end{cases}, \\ h(0,s') = {}& \frac{8}{27} - \frac{8}{9}\ln\frac{m_b}{\mu} - \frac{4}{9}\ln s' + \frac{4}{9}i\pi. \end{aligned} \tag{5}$$

Apart from this, the nonfactorizable effects [47–50] from the charm loop can bring about further corrections to the radiative $b \to s\gamma$ transition, which can be absorbed into the effective Wilson coefficient $C_7^{\rm eff}$. Specifically, the Wilson coefficient $C_7^{\rm eff}$ is given by [51]

$$C_7^{\rm eff}(\mu) = C_7(\mu) + C_{b\to s\gamma}(\mu),$$

with the absorptive part for the $b \to sc\bar{c} \to s\gamma$ rescattering

$$C_{b\to s\gamma}(\mu) = i\alpha_s\left[\frac{2}{9}\eta^{14/23}(G_1(x_t) - 0.1687) - 0.03C_2(\mu)\right], \tag{6}$$

$$G_1(x) = \frac{x(x^2-5x-2)}{8(x-1)^3} + \frac{3x^2\ln^2 x}{4(x-1)^4}, \tag{7}$$

where $\eta = \alpha_s(m_W)/\alpha_s(\mu)$, $x_t = m_t^2/m_W^2$. Here we have dropped out the tiny contributions proportional to CKM sector $V_{ub}V_{us}^*$. In addition, $C_7^{\prime\rm eff}(\mu)$ and $C_9^{\prime\rm eff}(\mu)$ can be obtained by replacing the unprimed Wilson coefficients with the corresponding prime ones in the above formulas.

III. PARAMETRIZATIONS OF MATRIX ELEMENTS AND FORM FACTORS IN LCSR

With the free quark decay amplitude available, we can proceed to calculate the decay amplitudes for semileptonic decays of $\bar{B}_0 \to K_0^*(1430)l^+l^-$ at hadronic level, which can be obtained by sandwiching the free quark amplitudes between the initial and final meson states. Consequently, the following two hadronic matrix elements

$$\begin{aligned} &\langle K_0^*(p)|\bar{s}\gamma_\mu\gamma_5 b|B_{q'}(p+q)\rangle, \\ &\langle K_0^*(p)|\bar{s}\sigma_{\mu\nu}\gamma_5 q^\nu b|B_{q'}(p+q)\rangle \end{aligned} \tag{8}$$

need to be computed as can be observed from Eq. (1). The contributions from vector and tensor types of transitions vanish due to parity conservations which is the property of strong interactions. Generally, the above two matrix elements can be parametrized in terms of a series of form factors as

$$\langle K_0^*(p)|\bar{s}\gamma_\mu\gamma_5 b|B_{q'}(p+q)\rangle = -i[f_+(q^2)p_\mu + f_-(q^2)q_\mu], \tag{9}$$

$$\langle K_0^*(p)|\bar{s}\sigma_{\mu\nu}\gamma_5 q^\nu b|B_{q'}(p+q)\rangle = -\frac{1}{m_B + m_{K_0^*}}[(2p+q)_\mu q^2 - (m_B^2 - m_S^2)q_\mu]f_T(q^2). \tag{10}$$

Contracting Eqs. (9) and (10) with the four momentum q^μ on both side and making use of the equations of motion

$$q^\mu(\bar{\psi}_1\gamma_\mu\psi_2) = (m_2 - m_1)\bar{\psi}_1\psi_2 \tag{11}$$

$$q^\mu(\bar{\psi}_1\gamma_\mu\gamma_5\psi_2) = -(m_1 + m_2)\bar{\psi}_1\gamma_5\psi_2 \tag{12}$$

we have

$$\langle K_0^*(p)|\bar{s}\gamma_5 b|B_{q'}(p+q)\rangle = \frac{-i}{m_b + m_s}[f_+(q^2)p\cdot q + f_-(q^2)q^2]. \tag{13}$$

To calculate the nonperturbative form factors, one has to rely on some nonperturbative approaches. Considering the distribution amplitudes up to twist-3, the form factors at small q^2 for $\bar{B}_0 \to K_0^* l^+l^-$ have been calculated in [36] using the LCSR. The dependence of form factors $f_i(q^2)$ $(i = +, -, T)$ on momentum transfer q^2 are parametrized

TABLE I. Numerical results for the parameters $f_i(0)$, a_i and b_i involved in the double-pole fit of form factors (15) responsible for $\bar{B}_0 \to K_0^*(1430) l \bar{l}$ decay up to the twist-3 distribution amplitudes of $K_0^*(1430)$ meson.

	$f_i(0)$	a_i	b_i
f_+	$0.97^{+0.20}_{-0.20}$	$0.86^{+0.19}_{-0.18}$	
f_-	$0.073^{+0.02}_{-0.02}$	$2.50^{+0.44}_{-0.47}$	$1.82^{+0.69}_{-0.76}$
f_T	$0.60^{+0.14}_{-0.13}$	$0.69^{+0.26}_{-0.27}$	

in either the single pole form

$$f_i(q^2) = \frac{f_i(0)}{1 - a_i q^2/m_{B_0}^2}, \tag{14}$$

or the double-pole form

$$f_i(q^2) = \frac{f_i(0)}{1 - a_i q^2/m_{B_0}^2 + b_i q^4/m_{B_{q_1}}^4}, \tag{15}$$

in the whole kinematical region $0 < q^2 < (m_{B_0} - m_{K_0^*})^2$ while nonperturbative parameters a_i and b_i can be fixed by the magnitudes of form factors corresponding to the small momentum transfer calculated in the LCSR approach. The results for the parameters a_i, b_i accounting for the q^2 dependence of form factors f_+, f_-, and f_T are grouped in Table I.

IV. FORMULA FOR PHYSICAL OBSERVABLES

In this section, we are going to perform the calculations of some interesting observables in phenomenology like the decay rates, forward-backward asymmetry as well as the polarization asymmetries of final state lepton. From Eq. (3), it is straightforward to obtain the decay amplitude for $\bar{B}_0 \to K_0^* l^+ l^-$ as

$$\mathcal{M}_{\bar{B}_0 \to K_0^* l^+ l^-} = -\frac{G_F \alpha}{2\sqrt{2}\pi} V_{tb} V_{ts}^* [T_\mu^1 (\bar{l}\gamma^\mu l) + T_\mu^2 (\bar{l}\gamma^\mu \gamma_5 l) + T^3 (\bar{l} l)], \tag{16}$$

where the functions T_μ^1, T_μ^2, and T^3 are given by

$$T_\mu^1 = i(C_9^{\rm eff} - C_9^{\prime \rm eff}) f_+(q^2) p_\mu + \frac{4 i m_b}{m_B + m_{K_0^*}} \times (C_7^{\rm eff} - C_7^{\prime \rm eff}) f_T(q^2) p_\mu, \tag{17}$$

$$T_\mu^2 = i(C_{10} - C_{10}')(f_+(q^2) p_\mu + f_-(q^2) q_\mu) - \frac{i}{2 m_l (m_b + m_s)} (C_{Q_2} - C_{Q_2}')(f_+(q^2) p \cdot q + f_-(q^2) q^2) q_\mu, \tag{18}$$

and

$$T^3 = i(C_{Q_1} - C_{Q_1}') \frac{1}{m_b + m_s} (f_+(q^2) p \cdot q + f_-(q^2) q^2). \tag{19}$$

It needs to point out that the terms proportional to q_μ in T_μ^1, namely $f_-(q^2)$ does not contribute to the decay amplitude with the help of the equation of motion for lepton fields. Besides, one can also find that the above results can indeed reproduce that obtained in the SM with $C_i' = 0$ and $T^3 = 0$.

A. The differential decay rates and forward-backward asymmetry of $\bar{B}_0 \to K_0^*(1430) l^+ l^-$

The semileptonic decay $\bar{B}_0 \to K_0^*(1430) l^+ l^-$ is induced by FCNCs. The differential decay width of $\bar{B}_0 \to K_0^*(1430) l^+ l^-$ in the rest frame of $\bar{B}_0$ meson can be written as [52]

$$\frac{d\Gamma(\bar{B}_0 \to K_0^*(1430) l^+ l^-)}{dq^2} = \frac{1}{(2\pi)^3} \frac{1}{32 m_{\bar{B}_0}} \int_{u_{\rm min}}^{u_{\rm max}} \times |\tilde{M}_{\bar{B}_0 \to K_0^*(1430) l^+ l^-}|^2 du, \tag{20}$$

where $u = (p_{K_0^*(1430)} + p_{l^-})^2$ and $q^2 = (p_{l^+} + p_{l^-})^2$; $p_{K_0^*(1430)}$, p_{l^+}, and p_{l^-} are the four-momenta vectors of $K_0^*(1430)$, l^+, and l^- respectively; $|\tilde{M}_{\bar{B}_0 \to K_0^*(1430) l^+ l^-}|^2$ is the squared decay amplitude after integrating over the angle between the lepton l^- and $K_0^*(1430)$ meson. The upper and lower limits of u are given by

$$u_{\rm max} = (E^*_{K_0^*(1430)} + E^*_{l^-})^2 - \left(\sqrt{E^{*2}_{K_0^*(1430)} - m^2_{K_0^*(1430)}} - \sqrt{E^{*2}_{l^-} - m^2_{l^-}}\right)^2,$$
$$u_{\rm min} = (E^*_{K_0^*(1430)} + E^*_{l^-})^2 - \left(\sqrt{E^{*2}_{K_0^*(1430)} - m^2_{K_0^*(1430)}} + \sqrt{E^{*2}_{l^-} - m^2_{l^-}}\right)^2, \tag{21}$$

where the energies of $K_0^*(1430)$ and l^- in the rest frame of lepton pair $E^*_{K_0^*(1430)}$ and $E^*_{l^-}$ are determined as

$$E^*_{K_0^*(1430)} = \frac{m_{\bar{B}_0}^2 - m^2_{K_0^*(1430)} - q^2}{2\sqrt{q^2}}, \qquad E_l^* = \frac{q^2}{2\sqrt{q^2}}. \tag{22}$$

Collecting everything together, one can write the general expression of the differential decay rate for $\bar{B}_0 \to K_0^*(1430) l^+ l^-$ as [53]:

$$\frac{d\Gamma}{ds} = \frac{G_F^2\alpha^2|V_{tb}V_{ts}^*|^2}{3072 m_B^3\pi^5 s}\sqrt{1-\frac{4m_l^2}{s}}\sqrt{\lambda(m_B^2, m_{K_0^*}^2, s)}\{|A|^2 \times (2m_l^2+s)\lambda + 12 s m_l^2(m_B^2 - m_{K_0^*}^2 - s)(CB^* + C^*B) + 12 m_l^2 s^2|C|^2 + 6t|D|^2(t-4m_l^2) + |B|^2((2m_l^2+s) \times (m_B^4 - 2m_B^2 m_{K_0^*}^2 - 2s m_{K_0^*}^2) + (m_{K_0^*}^2 - s)^2 + 2m_l^2(m_{K_0^*}^4 + 10 t m_{K_0^*}^2 + s^2))\}, \quad (23)$$

where

$$\lambda = \lambda(m_B^2, m_{K_0^*}^2, s) = m_B^4 + m_{K_0^*}^4 + s^2 - 2m_B^2 m_{K_0^*}^2 - 2m_{K_0^*}^2 s - 2s m_B^2. \quad (24)$$

The auxiliary functions are defined as

$$A = i(C_9^{\rm eff} - C_9^{\prime\rm eff})f_+(q^2) + \frac{4im_b}{m_B + m_{K_0^*}}(C_7^{\rm eff} - C_7^{\prime\rm eff})f_T(q^2)$$
$$B = i(C_{10} - C_{10}')f_+(q^2)$$
$$C = i(C_{10} - C_{10}')f_-(q^2) + \frac{i}{2m_e(m_b+m_s)}(p\cdot q f_+(q^2) + q^2 f_T(q^2))(C_{Q_2} - C_{Q_2}')$$
$$D = \frac{i}{m_b+m_s}(p\cdot q f_+(q^2) + q^2 f_T(q^2))(C_{Q_1} - C_{Q_1}') \quad (25)$$

The forward-backward asymmetry for the decay modes $\bar{B}_0 \to K_0^*(1430)l^+l^-$ is exactly equal to zero in the SM [54,55] due to the absence of scalar-type coupling between the lepton pair, which serves as a valuable ground to test the SM precisely as well as bound its extensions stringently. The differential forward-backward asymmetry of final state leptons in different SUSY models can be written as

$$\frac{dA_{\rm FB}(q^2)}{ds} = \int_0^1 d\cos\theta \frac{d^2\Gamma(s,\cos\theta)}{ds\,d\cos\theta} - \int_{-1}^0 d\cos\theta \frac{d^2\Gamma(s,\cos\theta)}{ds\,d\cos\theta} \quad (26)$$

and

$$A_{\rm FB}(q^2) = \frac{\int_0^1 d\cos\theta \frac{d^2\Gamma(s,\cos\theta)}{ds\,d\cos\theta} - \int_{-1}^0 d\cos\theta \frac{d^2\Gamma(s,\cos\theta)}{ds\,d\cos\theta}}{\int_0^1 d\cos\theta \frac{d^2\Gamma(s,\cos\theta)}{ds\,d\cos\theta} + \int_{-1}^0 d\cos\theta \frac{d^2\Gamma(s,\cos\theta)}{ds\,d\cos\theta}}. \quad (27)$$

Now putting everything together, we have

$$A_{\rm FB}(s) = \left(1/\frac{d\Gamma}{ds}\right)\frac{\alpha^2 G_F^2|V_{tb}V_{ts}^*|^2\lambda(m_B^2, m_{K_0^*}^2, s)}{1024 m_B^3\pi^5} \times m_l\left(1-\frac{4m_l^2}{s}\right)(AD^* + A^*D). \quad (28)$$

It is clear from the expressions of decay rate and forward-backward asymmetry that the contribution of the NHBs as well as that of the SUSY SO(10) GUT model comes in through the auxiliary functions defined in Eq. (25). Hence these SUSY effects manifest themselves in the numerical results of these observables.

B. Lepton polarization asymmetries of $\bar{B}_0 \to K_0^*(1430)l^+l^-$

In the rest frame of the lepton l^-, the unit vectors along longitudinal, normal and transversal component of the l^- can be defined as [56]:

$$s_L^{-\mu} = (0, \vec{e}_L) = \left(0, \frac{\vec{p}_-}{|\vec{p}_-|}\right),$$
$$s_N^{-\mu} = (0, \vec{e}_N) = \left(0, \frac{\vec{p}_{K_0^*}\times\vec{p}_-}{|\vec{p}_{K_0^*}\times\vec{p}_-|}\right), \quad (29)$$
$$s_T^{-\mu} = (0, \vec{e}_T) = (0, \vec{e}_N\times\vec{e}_L),$$

where $\vec{p}_-$ and $\vec{p}_{K_0^*}$ are the three-momenta of the lepton l^- and $K_0^*(1430)$ meson, respectively, in the center mass (CM) frame of l^+l^- system. Lorentz transformation is used to boost the longitudinal component of the lepton polarization to the CM frame of the lepton pair as

$$(s_L^{-\mu})_{\rm CM} = \left(\frac{|\vec{p}_-|}{m_l}, \frac{E_l\vec{p}_-}{m_l|\vec{p}_-|}\right), \quad (30)$$

where E_l and m_l are the energy and mass of the lepton. The normal and transverse components remain unchanged under the Lorentz boost.

The longitudinal (P_L), normal (P_N) and transverse (P_T) polarizations of lepton can be defined as:

$$P_i^{(\mp)}(s) = \frac{\frac{d\Gamma}{ds}(\vec{\xi}^{\mp} = \vec{e}^{\mp}) - \frac{d\Gamma}{ds}(\vec{\xi}^{\mp} = -\vec{e}^{\mp})}{\frac{d\Gamma}{ds}(\vec{\xi}^{\mp} = \vec{e}^{\mp}) + \frac{d\Gamma}{ds}(\vec{\xi}^{\mp} = -\vec{e}^{\mp})}, \quad (31)$$

where $i = L, N, T$, and $\vec{\xi}^{\mp}$ is the spin direction along the leptons $l^{\mp}$. The differential decay rate for polarized lepton $l^{\mp}$ in $\bar{B}_0 \to K_0^*(1430)l^+l^-$ decay along any spin direction $\vec{\xi}^{\mp}$ is related to the unpolarized decay rate (20) with the following relation

$$\frac{d\Gamma(\vec{\xi}^{\mp})}{ds} = \frac{1}{2}\left(\frac{d\Gamma}{ds}\right)[1 + (P_L^{\mp}\vec{e}_L^{\mp} + P_N^{\mp}\vec{e}_N^{\mp} + P_T^{\mp}\vec{e}_T^{\mp})\cdot\vec{\xi}^{\mp}]. \quad (32)$$

We can achieve the expressions of longitudinal, normal and transverse polarizations for $\bar{B}_0 \to K_0^*(1430)l^+l^-$ decays as collected below. The longitudinal lepton polarization can be written as

$$P_L(s) = \left(1/\frac{d\Gamma}{ds}\right)\frac{\alpha^2 G_F^2 |V_{tb}V_{ts}^*|^2 \lambda^{3/2}(m_B^2, m_{K_0^*}^2, s)}{3072 m_B^3 \pi^5} \times \left(1 - \frac{4m_l^2}{s}\right)(AB^* + A^*B). \quad (33)$$

Similarly, the normal lepton polarization is

$$P_N(s) = \left(1/\frac{d\Gamma}{ds}\right)\frac{\alpha^2 G_F^2 |V_{tb}V_{ts}^*|^2 m_l}{4096 m_B^3 \pi^4 \sqrt{s}}\sqrt{1 - \frac{4m_l^2}{s}} \times [(m_B^2 - m_{K_0^*}^2 + s)(A^*B + AB^*) - 2s(A^*C + AC^*)], \quad (34)$$

and the transverse one is given by

$$P_T(s) = \left(1/\frac{d\Gamma}{ds}\right)\frac{-i\alpha^2 G_F^2 |V_{tb}V_{ts}^*|^2 \lambda^{1/2}(m_B^2, m_{K_0^*}^2, s)}{4096 m_B^3 \pi^4} \times \left(1 - \frac{4m_l^2}{s}\right)(m_B^2 - m_S^2 + s) \times [(A^*D - AD^*) + 2m_l(B^*C - BC^*)]. \quad (35)$$

The $\frac{d\Gamma}{ds}$ appearing in the above equation is the one given in Eq. (23) and $\lambda(m_B^2, m_{K_0^*}^2, s)$ is the same as that defined in Eq. (24).

V. NUMERICAL ANALYSIS

In this section, we would like to present the numerical analysis of decay rates, forward-backward asymmetries and polarization asymmetries. The numerical values of Wilson coefficients and other input parameters used in our analysis are borrowed from Ref. [11,14,36] and collected in Tables II, III, and IV. In the subsequent analysis, we will focus on the parameter space of large $\tan\beta$, where the NHBs effects are significant owing to the fact that the Wilson coefficients corresponding to NHBs are proportional to $(m_b m_l / m_h^2)\tan^3\beta$ ($h = h^0, A^0$). Here, one $\tan\beta$ comes from the chargino-up-type squark loop and $\tan^2\beta$ comes from the exchange of the NHBs. At large value of $\tan\beta$ the $C_{Q_i}^{(\prime)}$ compete with $C_i^{(\prime)}$ and can overwhelm $C_i^{(\prime)}$ in some region as can be seen from the Tables III and IV [12]. SUSY I corresponds to the regions where SUSY can destructively contribute and can change the sign of C_7, but without contribution of NHBs, SUSY II refers to the region where $\tan\beta$ is large and the masses of the superpartners are relatively small. SUSY III corresponds to the regions where $\tan\beta$ is large and the masses of superpartners are relatively large. The primed Wilson coefficients are for the primed operators in Eq. (2) from NHBs contribution in SUSY SO(10) GUT model. As the NHBs are proportional to the lepton mass, the values shown in the table are for μ case and τ case (the values in parentheses of Table IV). Apart from the large $\tan\beta$ limit, the other two conditions responsible for the large contributions from NHBs are: (i) the mass values of the lighter chargino and lighter stop should not be too large; (ii) the mass splitting of charginos and stops should be large, which also indicate large mixing between stop sector and chargino sector [11]. Once these conditions are satisfied, the process $B \to X_s\gamma$ will not only impose constraints on C_7 but it also puts very stringent constraint on the possible new physics. It is well known that the SUSY contribution is sensitive to the sign of the Higgs mass term and SUSY contributes destructively when the sign of this term becomes minus. It is pointed out in literature [11] that there exist considerable regions of SUSY parameter space in which NHBs can largely contribute to the process $b \to s l^+ l^-$ due to change of the sign of C_7 from positive to negative, while the constraint on $b \to s\gamma$ is respected. Also, when the masses of SUSY particles are relatively large, say about 450 GeV, there exist significant regions in the parameter space of SUSY models in which NHBs could contribute largely. However, in these cases C_7 does not change its sign, because contributions of charged Higgs and charginos cancel each other.

Before, we discuss the numerical results, it is worth mentioning that the leptonic decay $B_q \to l\bar{l}$ is a golden channel to look for the NHBs effects in SUSY models at large $\tan\beta$. Its branching ratio can be written as [57–59]

TABLE II. Values of input parameters used in our numerical analysis

$G_F = 1.166 \times 10^{-5}$ GeV^{-2}	$\lvert V_{ts}\rvert = 41.61^{+0.10}_{-0.80} \times 10^{-3}$
$\lvert V_{tb}\rvert = 0.9991$	$m_b = (4.68 \pm 0.03)$ GeV
$m_c(m_c) = 1.275^{+0.015}_{-0.015}$ GeV	$m_s(1\text{ GeV}) = (142 \pm 28)$ MeV
$m_{B^0} = 5.28$ GeV	$m_{K_0^*} = 1.43$ GeV

TABLE III. Wilson Coefficients in SM and different SUSY models but without NHBs contributions. The primed Wilson coefficients corresponds to the operators which are opposite in helicities from those of the SM operators.

Wilson Coefficients	C_7^{eff}	$C_7^{\prime\text{eff}}$	C_9	C_9'	C_{10}	C_{10}'
SM	-0.313	0	4.334	0	-4.669	0
SUSY I or II	$+0.3756$	0	4.7674	0	-3.7354	0
SUSY III	-0.3756	0	4.7674	0	-3.7354	0
SUSY SO(10) ($A_0 = -1000$)	$-0.219 + 0i$	$0.039 - 0.038i$	$4.275 + 0i$	$0.011 + 0.0721i$	$-4.732 - 0i$	$-0.075 - 0.670i$

TABLE IV. Wilson coefficient corresponding to NHBs contributions. SUSY I corresponds to the regions where SUSY can destructively contribute and can change the sign of C_7, but without contribution of NHBs, SUSY II refers to the region where $\tan\beta$ is large and the masses of the superpartners are relatively small. SUSY III corresponds to the regions where $\tan\beta$ is large and the masses of superpartners are relatively large. The primed Wilson coefficients are for the primed operators in Eq. (2) from NHBs contribution in SUSY SO(10) GUT model. The values in the parentheses are for the τ case.

Wilson Coefficients	C_{Q_1}	C'_{Q_1}	C_{Q_2}	C'_{Q_2}
SM	0	0	0	0
SUSY I	0	0	0	0
SUSY II	6.5 (16.5)	0	-6.5 (-16.5)	0
SUSY III	1.2 (4.5)	0	-1.2 (-4.5)	0
SUSY SO(10) ($A_0 = -1000$)	$0.106 + 0i$	$-0.247 + 0.242i$	$-0.107 + 0i$	$-0.250 + 0.246i$
	$(1.775 + 0.002i)$	$(-4.148 + 4.074i)$	$(-1.797 - 0.002i)$	$(-4.202 + 4.128i)$

$$\mathrm{Br}(B_s \to \mu^+\mu^-) = \frac{G_F^2 \alpha_{\mathrm{em}}^2}{64\pi^3} m_{B_s}^3 \tau_{B_s} f_{B_s}^2 |V_{tb}^* V_{ts}|^2 \sqrt{1 - \frac{4m_l^2}{m_{B_s}^2}} \times \left[\left(1 - \frac{4m_l^2}{m_{B_s}^2}\right) C_{Q_1}^2 + \left(C_{Q_2} + \frac{2m_l}{m_{B_s}} C_{10}\right)^2\right], \quad (36)$$

where C_{Q_1} and C_{Q_2} corresponds to the NHB effects which are zero in the SM and are proportional to $(m_b m_l / m_h^2)\tan^3\beta$ ($h = h^0, A^0$) in MSSM. In SM, the value of Wilson coefficient C_{10} is large but this is suppressed by the factor $\frac{2m_l}{m_{B_s}}$ therefore, the expected value of the branching ratio is small and that is

$$\mathrm{BR}_{\mathrm{SM}}(B_s \to \mu^+\mu^-) = (3.86 \pm 0.15) \times 10^{-9}. \quad (37)$$

Now, with small and large value of the coefficients corresponding to the NHBs the branching ratio can be enhanced by a factor ranging from 10^1–10^3, but the current upper limit provided by CDF collaboration ($\mathrm{BR}_{\mathrm{SM}}(B_s \to \mu^+\mu^-) = 9.4 \times 10^{-8}$) [60] is just 25 times larger than the SM results. One can see clearly from Eq. (36) that the branching ratio for $B_s \to \mu^+\mu^-$ is directly proportional to the NHBs contributions. Therefore, this severely cuts the parameter space of the SUSY and the large value of the Wilson Coefficients C_{Q_1} and C_{Q_2} which are possible in SUSY II and SUSY III models have been ruled out. Only the value of these coefficients in SUSY SO(10) respect the current limit provided by CDF collaboration. Hence, the measurement of this decay at the future experiments will help to get useful constraints on the New Physics and if large deviation from SM predication is observed, then apart from the signal of supersymmetry, this would have an important implications on the Higgs searches at LHC.

The purpose of the study of NHBs effect in $B \to K_0^* l^+ l^-$ ($l = \mu, \tau$) is not only focused to incorporate the constraints provided by $B_s \to \mu^+\mu^-$ but to check the effects of NHBs at the larger extent. One can see that in our study we have chosen the value of C_{Q_1} and C_{Q_2} from 0 to 4.5 just to check the sensitivity of different physical observables on these NHBs. We hope that in future, when we have data on these decays, they will help us to test the constraints provided by the golden channel $B_s \to \mu^+\mu^-$.

The numerical results for the decay rates, forward-backward asymmetries and polarization asymmetries of the lepton are presented in Figs. 1–5. Figure 1 describes the differential decay rate of $B \to K^*(1430) l^+ l^-$, from

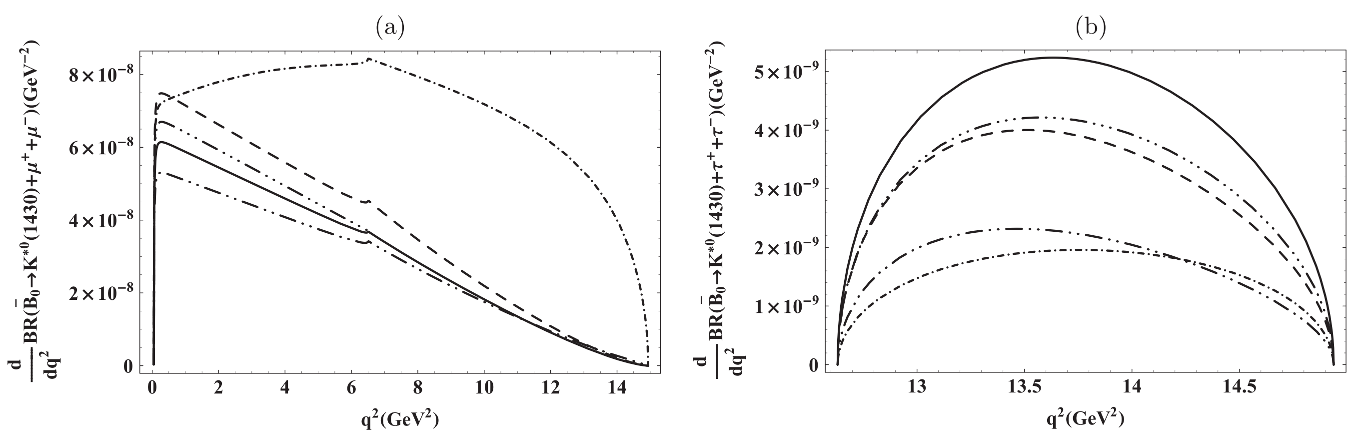

FIG. 1. The differential width for the $B \to K_0^* l^+ l^-$ ($l = \mu, \tau$) decays as functions of q^2. The solid, dashed, dashed-dot, dashed-double dot and dashed-triple dot line represents SM, SUSY I, SUSY II, SUSY III and SUSY SO(10) GUT model, respectively.

which one can see that the supersymmetric effects are quite distinctive from that of the SM both in the small and large momentum region. The reason for the increase of differential decay width in SUSY I model is the relative change in the sign of $C_7^{\rm eff}$; while the large change in SUSY II model is due to the contribution of the NHBs. As for the SUSY III and SUSY SO (10) models, the value of the Wilson coefficients corresponding to NHBs is small and hence one expects small deviations from SM. For SUSY II, the value of decay rate is smaller than that of SM value which is due to the change in the sign of $C_7^{\rm eff}$ and also due to the small contributions from the NHB's. Similar effects can also be seen for the tauon case in Fig. 1(b). For this case, the SUSY play destructive role which is quite different compared to $B \to K l^+ l^-$ decays.

As an exclusive decay, there are different source of uncertainties involved in the calculation of the above said decay. The major uncertainties in the numerical analysis of $\bar{B}_0 \to K_0^*(1430) l^+ l^-$ decay originated from the $\bar{B}_0 \to K_0^*(1430)$ transition form factors calculated in the LCSR approach as shown in Table I, which can bring about almost 40% errors to the differential decay rate of $\bar{B}_0 \to K_0^*(1430) l^+ l^-$, which showed that it is not a suitable tool to look for the new physics. The large uncertainties involved in the form factors are mainly from the variations of the decay constant of $K_0^*(1430)$ meson and the Gengenbauer moments in its distribution amplitudes. There are also some uncertainties from the strange quark mass m_s, which are expected to be very tiny on account of the negligible role of m_s suppressed by the much larger energy scale of m_b. Moreover, the uncertainties of the charm quark and bottom quark mass are at the 1% level, which will not play significant role in the numerical analysis and can be dropped out safely. It also needs to be stressed that these hadronic uncertainties almost have no influence on the various asymmetries including the lepton polarization asymmetry on account of the serious cancellation among different polarization states and this make them the best tool to look for physics beyond the SM.

In Fig. 2, the forward-backward asymmetries for $B \to K_0^* l^+ l^-$ ($l = \mu$, τ) are presented. In SM the forward-backward asymmetry is zero for this decay because there is no scalar operator. However in SUSY II, SUSY III and SUSY SO(10) model, we have the scalar operators corresponding to the NHBs, therefore we expect the nonzero value of the forward-backward asymmetry. This is quite clear from the Eq. (28) where the auxiliary function D corresponds to the contributions from NHBs. Figure 2(a) describes the forward-backward asymmetry for $B \to K_0^* \mu^+ \mu^-$. As the forward-backward asymmetry is proportional to the lepton mass, therefore for the muons case it is expected to be very small compared to the tauons case. Thus the maximum value of the forward-backward asymmetry is 0.05 in SUSY II model which is hard to be observed experimentally. However, for $B \to K_0^* \tau^+ \tau^-$ the maximum value of forward-backward asymmetry is around 0.35 in SUSY II model. Now this value differs from $B \to K l^+ l^-$ only in the sign [11], and hence the experimental results of this decay will help us in understanding the SUSY effects in B meson decays. The number of events required to observe this asymmetry are around 10^8 or so which are accessible at large colliders like the LHCb. When the final state leptons are the tauon pair, the effects of SUSY III and SUSY SO(10) are still too small to be measured experimentally.

Figure 3(a) and 3(b) shows the dependence of longitudinal polarization asymmetry for the $B \to K_0^* l^+ l^-$ on the square of momentum transfer. The value of longitudinal lepton polarization for muon is around -1 in the SM and we have a slight deviation on this value for SUSY I and SUSY SO(10) model. However, in SUSY II and SUSY III model the value of longitudinal lepton polarization approaches to zero in the large momentum transfer region. The reason is that in SUSY II model we have a large value

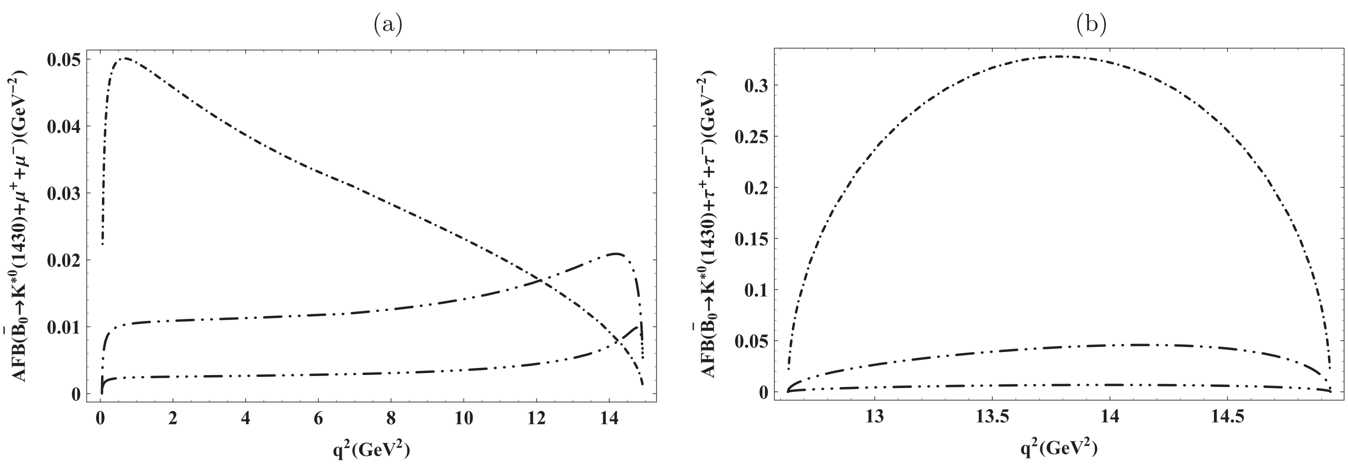

FIG. 2. Forward-backward asymmetry for the $B \to K_0^* l^+ l^-$ ($l = \mu$, τ) decays as functions of q^2. The dashed-dot, dashed-double dot and dashed-triple dot line represents SUSY II, SUSY III and SUSY SO(10) GUT model, respectively.

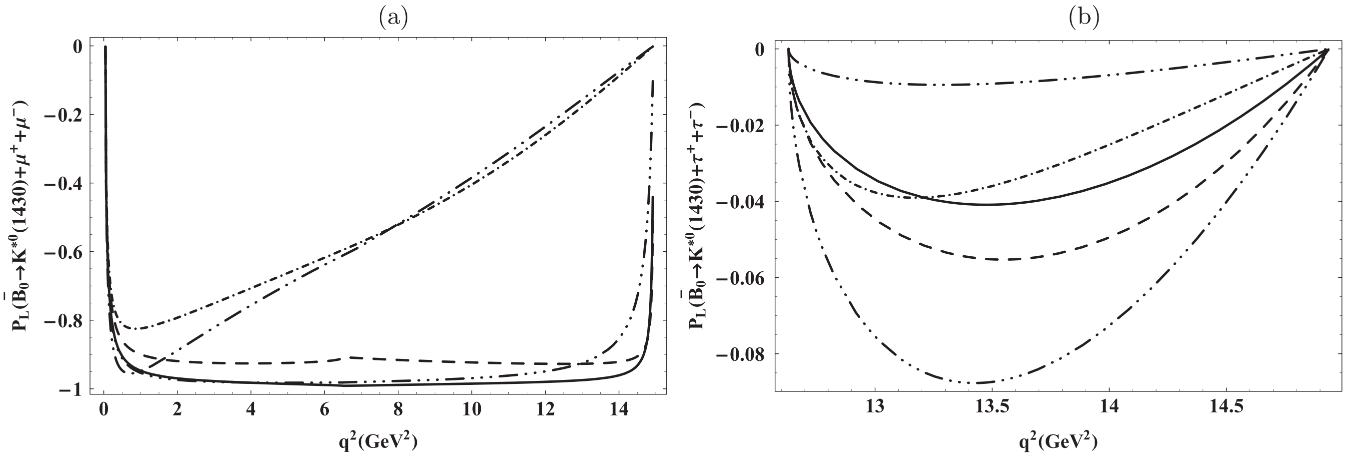

FIG. 3. Longitudinal lepton polarization asymmetries for the $B \to K_0^* l^+ l^-$ ($l = \mu, \tau$) decays as functions of q^2. The solid, dashed, dashed-dot, dashed-double dot and dashed-triple dot line represents SM, SUSY I, SUSY II, SUSY III and SUSY SO(10) GUT model, respectively.

of the differential decay rate and this suppresses the value of the polarization in the large q^2 region. In SUSY III though the value of the decay rate is not large, relatively small contribution comes from the Wilson coefficients $C_7^{*\text{eff}} C_{10}$. In large q^2 region, the longitudinal lepton polarization approaches to zero in all the models including the SM, because the factor $\lambda(m_B^2, m_{K_0^*}^2, s)$ goes to zero at large value of transfer momentum. Similar effects can be seen for the final state tauon but the value for this case is too small to measure experimentally.

The dependence of lepton normal polarization asymmetries for $B \to K_0^* l^+ l^-$ on the momentum transfer squared are presented in Fig. 4. In terms of Eq. (34), one can observe that this asymmetry is sensitive to the contribution of NHBs in almost all the supersymmetric models. Figure 4(a) shows the normal lepton polarization for $B \to K_0^* \mu^+ \mu^-$. It can be seen that P_N changes its sign in the case of SUSY III model and this is due to the contribution from NHBs. Now for SUSY II model, though large contributions from NHBs but it is overshadowed by the opposite sign of $C_7^{*\text{eff}}$ and C_9^{eff}. As the normal lepton polarization is proportional to the lepton mass, for $\tau^+\tau^-$ channel, it is expected that one can distinguish between different SUSY models, which can be seen from the Fig. 4(b). Again due to same reasons as for the muons case, the normal lepton polarization changes its sign in SUSY III model. The main advantage to study $B \to K_0^* l^+ l^-$ decay is that the value of normal lepton polarization is very large compared to $B \to K(K^*) l^+ l^-$ studied in these model in Ref. [11] and the reason for this is quite obvious. From Eq. (34) one can see that the normal lepton polarization for $B \to K_0^* l^+ l^-$ is proportional to the real part of the Wilson coefficients which is contrary to the $B \to K(K^*) l^+ l^-$ decays where it is proportional to the imaginary part [11]. As the imaginary part of these coefficients is very small in SUSY I, SUSY II and SUSY III models, therefore the corresponding value of the normal polarization is also small. This really makes it quite interesting to

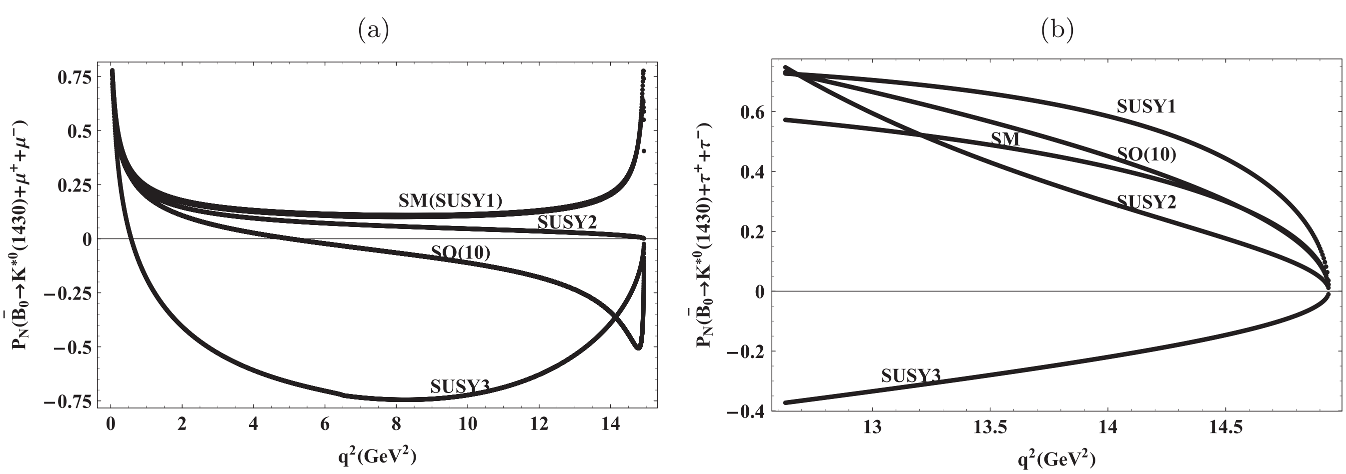

FIG. 4. Normal lepton polarization asymmetries for the $B \to K_0^* l^+ l^-$ ($l = \mu, \tau$) decays as functions of q^2.

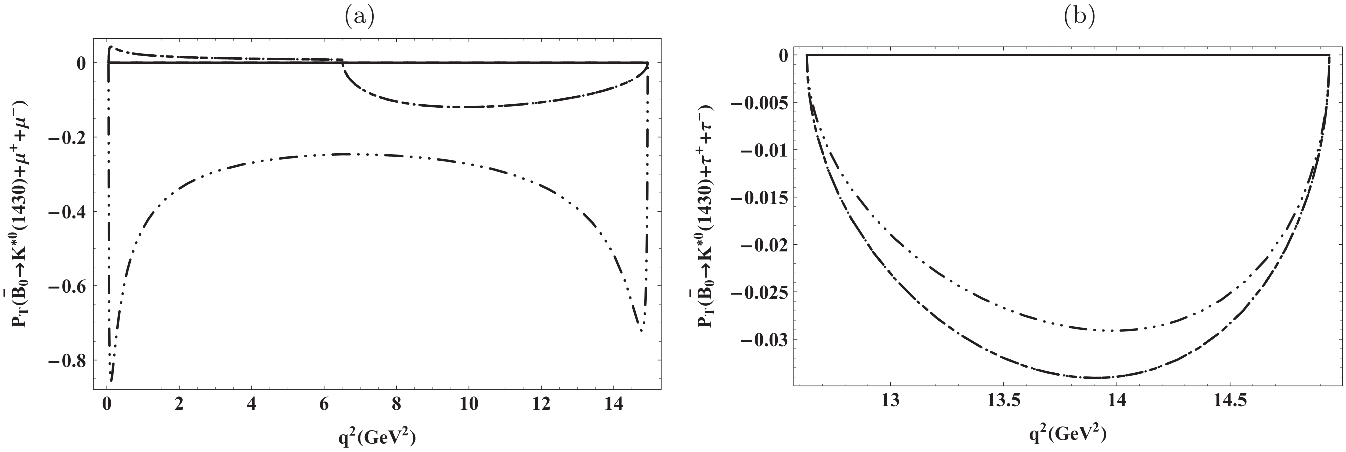

FIG. 5. Transverse lepton polarization asymmetries for the $B \to K_0^* l^+ l^-$ ($l = \mu, \tau$) decays as functions of q^2.

look for SUSY signatures in $B \to K_0^* l^+ l^-$ decays in the future experiments.

Figure 5 shows the dependence of transverse polarization asymmetries for $B \to K_0^* l^+ l^-$ on the square of momentum transfer. From Eq. (35) we can see that it is proportional to the imaginary part of the Wilson coefficients which are negligibly small in SM as well as in SUSY I, SUSY II and SUSY III models. However, complex flavor nondiagonal down-type squark mass matrix elements of 2nd and 3rd generations are of order one at GUT scale in SUSY SO(10) model, which induce complex couplings and Wilson coefficients. As a result, non zero transverse polarization asymmetries for $B \to K_0^* l^+ l^-$ exist in this model. Now for $\mu^+\mu^-$ channel, the value of transverse polarization asymmetry is around -0.3 in almost all value of q^2 except at the end points.

Finally, we take the transverse polarization asymmetry $\langle P_T \rangle$ as an example to show the discovery potential of SUSY with the future experiments. Experimentally, to measure $\langle P_T \rangle$ of a particular decay branching ratio $\mathcal{B}$ at the $n\sigma$ level, the required number of events are $N = n^2/(\mathcal{B}\langle P_T \rangle^2)$ and if $\langle P_T \rangle \sim 0.3$, then the required number of events are almost 10^8 for B decays. Since at the large hadron collider, the expected number of $b\bar{b}$ production events is around 10^{12} per year, so the measurement of transverse polarization asymmetries in the $B \to K_0^* l^+ l^-$ decays could discriminate the SUSY SO(10) model from the SM and other SUSY models. SuperB factory, having an initial luminosity of 10^{36} cm^{-2} s^{-1} can collect $15ab^{-1}$ data samples in a New Snowmass Year, which will also open up the world of New Physics effects including SUSY in such rare B decays.

VI. CONCLUSION

We have carried out the study of invariant mass spectrum, forward-backward asymmetry, polarization asymmetries of semileptonic decays $B \to K_0^*(1430)l^+l^-$ ($l = \mu, \tau$) in SUSY theories including SUSY SO(10) GUT model. Particularly, we analyzed the effects of NHBs to this process and the main outcomes of this study can be summarized as follows:

(i) The differential decay rates deviate sizably from that of the SM especially in the large momentum transfer region. These effects are significant in SUSY II model where the value of the Wilson coefficients corresponding to the NHBs is large. However, the SUSY SO(10) effects in differential decay rate of $B \to K_0^*(1430)l^+l^-$ ($l = \mu, \tau$) are negligibly small.

(ii) The forward-backward asymmetry for the decay $B \to K_0^* l^+ l^-$ is zero in the SM because of the missing of scalar operators in SM. Hence, the SUSY effects show up and the maximum value of the forward-backward asymmetry is around 0.35 for $B \to K_0^* \tau^+ \tau^-$ in SUSY II model. When the final state leptons are the tauon pair, the effects of SUSY III and SUSY SO(10) are still too small to be measured experimentally.

(iii) The longitudinal, normal and transverse polarizations of leptons are calculated in different SUSY models. It is found that the SUSY effects are very promising which could be measured at future experiments and shed light on the new physics signal beyond the SM. The transverse polarization asymmetry is the most interesting observable to look for the SUSY SO(10) effects where its value is around 0.3 in almost all the q^2 region. It is measurable at future experiments like LHC and BTeV machines where a large number of $b\bar{b}$ pairs are expected to be produced.

In short, the experimental investigation of observables, like decay rates, forward-backward asymmetry and lepton polarization asymmetries in $B \to K_0^*(1430)l^+l^-$ ($l = \mu, \tau$) decay will be used to search for the SUSY effects, in particular, the NHBs effect, encoded in the MSSM as well as SUSY SO(10) models.

ACKNOWLEDGMENTS

This work is partly supported by National Science Foundation of China under Grant No. 10735080 and 10625525. We acknowledge Professor Riazuddin for the helpful discussions and Jamil Aslam would like to thank C. D. Lü for the kind hospitality in Institute of High Energy Physics, CAS, where part of this work has been done.

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