The weak decays of $B \to K^*_0(1430)l^+l^-$ are investigated in Minimal Supersymmetric Standard Model (MSSM) and also in Supersymmetric (SUSY) SO(10) Grand Unified 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) GUT 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 (SM) 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.

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 Cabibbo-Kobayashi-Maskawa (CKM) matrix elements as well as the various hadronic form factors. In literature, there have been intensive studies on the exclusive decays $B \to P(V,A) l^+l^-$ [2, 3, 4, 5, 6, 7, 8, 9, 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 SM no one can say that it is the ultimate theory of nature as it has many open questions, such as 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 \( B \rightarrow X_s l^+l^- \), 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 \( B \rightarrow K(K^*) l^+l^- \) 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 Models (GUT) [13] has been proposed in literature. In this model, there is a complex flavor non-diagonal 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 \( B \rightarrow K(K^*) l^+l^- \) 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].

In this paper, we will investigate the exclusive decay \( \bar{B}_0 \rightarrow K^*_0(1430) \bar{u}l (l = \mu, \tau) \), 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]. 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 [13]. 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 \( B \rightarrow X_s \gamma \) 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 semi-leptonic \( B \) meson decay is to properly evaluate the hadronic matrix elements for \( B \rightarrow K^*_0(1430) \), namely the transition form factors, which are governed by the non-perturbative 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 non-perturbative 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, 22, 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 non-perturbative 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 semi-leptonic decays of heavy hadrons [25, 26, 27, 28], radiative hadronic decays [29, 30, 31] and non-leptonic two body decays of $B$ meson [32, 33, 34, 35].

In our numerical analysis for $\bar{B}_0 \rightarrow K^*(1430)$ decays, we shall use the results of the form factors calculated by LCSR approach in Ref. [36], and 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. 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.

The paper is organized as follows. In Sec. II, we present the effective Hamiltonian for the semileptonic decay $B \rightarrow K^* l^+ l^-$. Section III contains the parameterizations 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 $b \rightarrow s l^+ l^-$ in SUSY SO(10) GUT model, can be written as [14]

$$H_{\text{eff}} = -\frac{4G_F}{\sqrt{2}} \sum_{i=1}^{2} C_i(\mu) O_i(\mu) + \sum_{i=3}^{10} (C_i(\mu) O_i(\mu) + C'_i(\mu) O'_i(\mu)) + \sum_{i=1}^{8} (C_Q_i(\mu) Q_i(\mu) + C'_Q_i(\mu) Q'_i(\mu)),$$

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 $\mu$. 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 chromo-magnetic 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 $\delta_{23}^{RR}$ can be as large as 0.5 [13, 14]. Apart from this, $\delta_{23}^{RR}$ 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 $B \rightarrow K_0^*(1430)l^+l^-$ transition are given by

$$O_7 = \frac{e^2}{16\pi^2}m_b (\bar{s}\sigma_{\mu\nu}P_Rb) F^{\mu\nu}, \quad O_7' = \frac{e^2}{16\pi^2}m_b (\bar{s}\sigma_{\mu\nu}P_Lb) F^{\mu\nu}$$

$$O_9 = \frac{e^2}{16\pi^2}(\bar{s}\gamma_\mu P_Lb)(\bar{l}\gamma^\mu l), \quad O_9' = \frac{e^2}{16\pi^2}(\bar{s}\gamma_\mu P_Rb)(\bar{l}\gamma^\mu l)$$

$$O_{10} = \frac{e^2}{16\pi^2}(\bar{s}\gamma_\mu P_Lb)(\bar{l}\gamma^\mu s\gamma_5 l), \quad O_{10}' = \frac{e^2}{16\pi^2}(\bar{s}\gamma_\mu P_Rb)(\bar{l}\gamma^\mu s\gamma_5 l)$$

$$Q_1 = \frac{e^2}{16\pi^2}(\bar{s}P_Rb)(\bar{l}l), \quad Q_1' = \frac{e^2}{16\pi^2}(\bar{s}P_Lb)(\bar{l}l)$$

$$Q_2 = \frac{e^2}{16\pi^2}(\bar{s}P_Rb)(\bar{l}\gamma_5 l), \quad Q_2' = \frac{e^2}{16\pi^2}(\bar{s}P_Lb)(\bar{l}\gamma_5 l)$$

(2)

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

$$\mathcal{M}(b \rightarrow s l^+l^-) = -\frac{G_F\alpha}{\sqrt{2\pi}}V_{tb}V_{ts}^* \left\{ C_7^{eff}(\bar{s}\gamma_\mu P_Lb)(\bar{l}\gamma^\mu l) + C_10(\bar{s}\gamma_\mu P_Lb)(\bar{l}\gamma^\mu s\gamma_5 l) -2m_b C_7^{eff}(\bar{s}\sigma_{\mu\nu}\frac{q^\nu}{8} P_Rb)(\bar{l}\gamma^\mu l) + C_1 Q_1(\bar{s}P_Rb)(\bar{l}l) + C_2 Q_2(\bar{s}P_Rb)(\bar{l}\gamma_5 l) + (C_i(m_b) \leftrightarrow C'_i(m_b)) \right\},$$

(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, $\sum_{i=1}^6 (l^+l^-s|O_i|b)$, which are usually absorbed into the effective Wilson coefficient $C_9^{eff}(\mu)$, that one can decompose into the following three parts [40, 41, 42, 43, 44, 45, 46, 47, 48, 49, 50]

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

where the parameters $z$ and $s'$ are defined as $z = m_c/m_b$, $s' = q^2/m_b^2$. $Y_{SD}(z, s')$ 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 $Y_{LD}(z, s')$ from four-quark operators near the $c\bar{c}$ resonance cannot be calculated from first principles of QCD and are usually parameterized 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 $B \rightarrow J/\psi K_0^*(1430)$. The
manifest expressions for $Y_{SD}(z, s')$ can be written as

$$Y_{SD}(z, s') = h(z, s')(3C_1(\mu) + C_2(\mu) + 3C_3(\mu) + C_4(\mu) + 3C_5(\mu) + C_6(\mu))$$

$$-\frac{1}{2}h(1, s')(4C_3(\mu) + 4C_4(\mu) + 3C_5(\mu) + C_6(\mu))$$

$$-\frac{1}{2}h(0, s')(C_3(\mu) + 3C_4(\mu)) + \frac{2}{9}(3C_3(\mu) + C_4(\mu) + 3C_5(\mu) + C_6(\mu)),$$

(4)

with

$$h(z, s') = -\frac{8}{9}\ln z + \frac{8}{27} + \frac{4}{9}x - \frac{2}{9}(2 + x)|1 - x|^{1/2} \left\{ \ln \left(\frac{\sqrt{x - 1} + 1}{\sqrt{x - 1} - 1}\right) - i\pi \right\} \quad \text{for } x \equiv 4z^2/s' < 1,$$

$$2 \arctan \frac{1}{\sqrt{x - 1}} \quad \text{for } x \equiv 4z^2/s' > 1,$$

(5)

Apart from this, the non-factorizable effects 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 $C_7^{\text{eff}}$. Specifically, the Wilson coefficient $C_7^{\text{eff}}$ is given by

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

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

$$C_{b \to s\gamma}(\mu) = i\alpha_s \left[ \frac{2}{3} \eta^{14/23}(G_1(x_1) - 0.1687) - 0.03C_2(\mu) \right],$$

(6)

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

(7)

where $\eta = \frac{\alpha_s(m_W)}{\alpha_s(\mu)}$, $x_1 = m_t^2/m_W^2$. Here we have dropped out the tiny contributions proportional to CKM sector $V_{ub}V_{us}^\ast$. In addition, $C_7^{\text{eff}}(\mu)$ and $C_9^{\text{eff}}(\mu)$ can be obtained by replacing the unprimed Wilson coefficients with the corresponding prime ones in the above formulæ.

### III. PARAMETERIZATIONS OF MATRIX ELEMENTS AND FORM FACTORS IN LCSR

With the free quark decay amplitude available, we can proceed to calculate the decay amplitudes for semi-leptonic decays of $B_0 \to K_0^*(1430)^{+}\pi^-$ 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

$$\langle K_0^*(p)|\bar{s}\gamma_\mu\gamma_5b|B_0^\prime(p + q)\rangle, \quad \langle K_0^*(p)|\bar{s}\sigma_{\mu\nu}\gamma_5q^\nu b|B_0^\prime(p + q)\rangle$$

(8)

need to be computed as can be observed from Eq. 11. 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 parameterized in terms of a series of form factors as

$$\langle K_0^*(p)|\bar{s}\gamma_\mu\gamma_5b|B_0^\prime(p + q)\rangle = -i[f_+(q^2)p_\mu + f_-(q^2)q_\mu],$$

(9)

$$\langle K_0^*(p)|\bar{s}\sigma_{\mu\nu}\gamma_5q^\nu b|B_0^\prime(p + q)\rangle = -\frac{1}{m_B + m_{K_0^*}}\left(2p + q\right)_\mu \left[q^2 - (m_B^2 - m_{K_0^*}^2)q_\mu\right] f_T(q^2).$$

(10)
At the large recoil region, these form factors satisfy the following relations \[36, 52\]
\[
f_+(q^2) = \frac{2m_B}{m_B + m_{K^*_0}} f_T(q^2), \quad f_-(q^2) = 0, \tag{11}
\]
\[
f_T(q^2) = -\frac{m_b - m_{K^*_0}}{m_B - m_{K^*_0}} f_+(q^2), \tag{12}
\]
with the help of the equations of motion in the heavy quark limit \[53\]. Contracting Eqs. \[9,10\] with the four momentum $q^\mu$ on both side and making use of the equations of motion
\[
q^\mu (\bar{\psi}_1 \gamma_\mu \psi_2) = (m_2 - m_1) \bar{\psi}_1 \psi_2 \tag{13}
\]
\[
q^\mu (\bar{\psi}_1 \gamma_\mu \gamma_5 \psi_2) = -(m_1 + m_2) \bar{\psi}_1 \gamma_5 \psi_2 \tag{14}
\]
we have
\[
\langle K^*_0(p) | \bar{s} \gamma_5 b | B_q(p + q) \rangle = \frac{-i}{m_b + m_s} [f_+(q^2)p \cdot q + f_-(q^2)q^2] \tag{15}
\]

To calculate the non-perturbative 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 parameterized in either the single pole form
\[
f_i(q^2) = \frac{f_i(0)}{1 - a_i q^2 / m_{B_0}^2}, \tag{16}
\]
or the double-pole form
\[
f_i(q^2) = \frac{f_i(0)}{1 - a_i q^2 / m_{B_0}^2 + b_i q^4 / m_{B_{41}}^4}, \tag{17}
\]
in the whole kinematical region $0 < q^2 < (m_{B_0} - m_{K^*_0})^2$ while non-perturbative 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 \[I\].

**TABLE I**: Numerical results for the parameters $f_i(0)$, $a_i$ and $b_i$ involved in the double-pole fit of form factors \[17\] 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.

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

**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. \[43\],
it is straightforward to obtain the decay amplitude for $\bar{B}_0 \to K^*_0(1430)^+ l^-$ as

$$
M_{\bar{B}_0 \to K^*_0(1430)^+ l^-} = \frac{G_F \alpha}{2\sqrt{2\pi}} V_{tb} V^*_{ts} \left[ T^1_{\mu} (i\gamma^\mu l) + T^2_{\mu} (i\gamma^\mu \gamma_5 l) + T^3 (\bar{l}l) \right],
$$

where the functions $T^1_{\mu}$, $T^2_{\mu}$ and $T^3$ are given by

$$
T^1_{\mu} = i \left( C_{\bar{g}} - C_{q} \right) f_+(q^2)p_{\mu} + \frac{4i m_b}{m_B + m_{K^*}} \left( C_{\bar{g}} - C_{q} \right) f_+(q^2)p_{\mu},
$$

$$
T^2_{\mu} = i \left( C_{10} - C'_{10} \right) (f_+(q^2)p_{\mu} + f_-(q^2)q_{\mu}) - \frac{i}{2 m_t (m_b + m_s)} \left( C_{Q_2} - C'_{Q_2} \right) (f_+(q^2)p \cdot q + f_-(q^2)q) q_{\mu},
$$

and

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

It needs to point out that the terms proportional to $q_{\mu}$ in $T^3_{\mu}$, 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 $\bar{B}_0 \to K^*_0(1430)^+ l^-$

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

$$
\frac{d\Gamma(\bar{B}_0 \to K^*_0(1430)^+ l^-)}{d q^2} = \frac{1}{16\pi} \frac{1}{32 m_{B_0}} \int_{u_{\text{min}}}^{u_{\text{max}}} |\tilde{M}_{\bar{B}_0 \to K^*_0(1430)^+ l^-}|^2 du,
$$

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; $|\tilde{M}_{\bar{B}_0 \to K^*_0(1430)^+ l^-}|^2$ 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

$$
u_{\text{max}} = (E^*_{K^*_0(1430)} + E^*_{l^-})^2 - \sqrt{(E^*_{K^*_0(1430)}^2 - m_{K^*_0(1430)}^2 - m_{l^-}^2)},
$$

$$
u_{\text{min}} = (E^*_{K^*_0(1430)} + E^*_{l^-})^2 - \sqrt{(E^*_{K^*_0(1430)}^2 - m_{K^*_0(1430)}^2 + m_{l^-}^2)},
$$

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

$$
E^*_{K^*_0(1430)} = \frac{m_{B_0}^2 - m_{K^*_0(1430)}^2 - q^2}{2\sqrt{q^2}}, \quad E^*_{l^-} = \frac{q^2}{2\sqrt{q^2}}.
$$

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

$$
\frac{d\Gamma}{ds} = \frac{G_F^2 \alpha^2 |V_{tb} V^*_{ts}|^2}{3072 m_{B_0} \pi^3 s} \sqrt{1 - \frac{4m_t^2}{s}} \sqrt{x m_{B_0}^2 m_{K^*_0}, s} \times 
\left\{ |A|^2 \left( 2m_l^2 + s \right) \lambda + 12s \left( m_B^2 - m_{K^*_0}^2 - s \right) \left( CB^* + C^*B \right) + 12m_t^2 s^2 |C|^2 + 6t |D|^2 (t - 4m_t^2) \right\},
$$

$$
+ |B|^2 \left( 2m_l^2 + s \right) \left( m_B^4 - 2m_B^2 m^2_{K^*_0} - 2sm^2_{K^*_0} \right) + \left( m_{K^*_0}^2 - s \right)^2 + 2m_t^2 \left( m_{K^*_0}^2 + 10tm_{K^*_0}^2 + s^2 \right),
$$

(25)
where

$$\lambda = \lambda(m_B^2, m_{K_0}^2, s) = m_B^4 + m_{K_0}^4 + s^2 - 2m_B^2m_{K_0}^2 - 2m_{K_0}^2s - 2m_B^2. \quad (26)$$

The auxiliary functions are defined as

$$A = i\left(C^{\text{eff}}_6 - C^{\text{eff}}_7\right) f_+(q^2) + \frac{4im_b}{m_B + m_{K_0}}\left(C^{\text{eff}}_6 - C^{\text{eff}}_7\right) f_T(q^2)$$

$$B = i(C_{10} - C'_{10}) f_+(q^2)$$

$$C = i(C_{10} - C'_{10}) f_-(q^2) + \frac{i}{2m_e(m_b + m_s)}(p \cdot q f_+(q^2) + q^2 f_T(q^2))(C_{Q_2} - C'_{Q_2})$$

$$D = \frac{i}{m_b + m_s}(p \cdot q f_+(q^2) + q^2 f_T(q^2))(C_{Q_1} - C'_{Q_1}) \quad (27)$$

The forward-backward asymmetry for the decay modes $\bar{B}_0 \rightarrow K_0^*(1430)t^+l^-$ is exactly equal to zero in the SM 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

$$\frac{dA_{FB}(q^2)}{ds} = \int_0^1 d\cos\theta \frac{d^2\Gamma(s, \cos\theta)}{dsd\cos\theta} - \int_{-1}^0 d\cos\theta \frac{d^2\Gamma(s, \cos\theta)}{dsd\cos\theta} \quad (28)$$

and

$$A_{FB}(q^2) = \int_0^1 d\cos\theta \frac{d^2\Gamma(s, \cos\theta)}{dsd\cos\theta} - \int_{-1}^0 d\cos\theta \frac{d^2\Gamma(s, \cos\theta)}{dsd\cos\theta} \quad (29)$$

Now putting everything together, we have

$$A_{FB}(s) = \left(1 + \frac{G_F^2|V_{tb}V_{ts}^*|^2 \lambda(m_B^2, m_{K_0}^2, s)}{1024m_B^3\pi^5} m_l(1 - \frac{4m_l^2}{s})\right)(AD^* + A^*D). \quad (30)$$

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. (27). Hence these SUSY effects manifest themselves in the numerical results of these observables.

### B. Lepton Polarization asymmetries of $\bar{B}_0 \rightarrow K_0^*(1430)t^+l^-$

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

$$s_L^{-\mu} = (0, \hat{e}_L) = \left(0, \frac{\vec{p}_-}{|\vec{p}_-|}\right),$$

$$s_N^{-\mu} = (0, \hat{e}_N) = \left(0, \frac{\vec{p}_{K_0} \times \vec{p}_-}{|\vec{p}_{K_0} \times \vec{p}_-|}\right),$$

$$s_T^{-\mu} = (0, \hat{e}_T) = \left(0, \hat{e}_N \times \hat{e}_L\right), \quad (31)$$

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 $t^+l^-$ system. Lorentz transformation is used to boost the longitudinal component of the lepton polarization to the CM frame of the lepton pair as

$$(s_L^{-\mu})_{CM} = \left(\frac{|\vec{p}_-|}{m_l}, \frac{E_l|\vec{p}_-|}{m_l|\vec{p}_-|}\right) \quad (32)$$
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:

$$P^\pm_i(s) = \frac{d\Gamma(\xi^\pm)}{ds}(\xi^\pm = \epsilon^\pm) - \frac{d\Gamma(\xi^\pm)}{ds}(\xi^\pm = -\epsilon^\pm)$$

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

$$\frac{d\Gamma(\xi^\pm)}{ds} = \frac{1}{2} \left( \frac{d\Gamma}{ds} \right) [1 + (P_L^\pm \epsilon_L^\pm + P_N^\pm \epsilon_N^\pm + P_T^\pm \epsilon_T^\pm) \cdot \xi^\pm].$$

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

$$P_L(s) = (1/d\Gamma) \frac{\alpha^2 G_F^2 |V_{tb}V_{ts}^*|^2 \lambda^{3/2}(m^2_B, m^2_{K^*_0}, s)}{3072m^4_B\pi^5} \left( 1 - \frac{4m_l^2}{s} \right) (A^*B + A^*B).$$

Similarly, the normal lepton polarization is

$$P_N(s) = (1/d\Gamma) \frac{\alpha^2 G_F^2 |V_{tb}V_{ts}^*|^2 m_l}{4096m^3_B\pi^4\sqrt{s}} \left[ 1 - \frac{4m_l^2}{s} \right] \left( m^2_B - m^2_{K^*_0} + s \right) (A^*B + A^*B) - 2s(A^*C + AC^*).$$

and the transverse one is given by

$$P_T(s) = \left( 1/d\Gamma \right) \frac{-i\alpha^2 G_F^2 |V_{tb}V_{ts}^*|^2 \lambda^{1/2}(m^2_B, m^2_{K^*_0}, s)}{4096m^4_B\pi^4} \left[ 1 - \frac{4m_l^2}{s} \right] \left( m^2_B - m^2_{K^*_0} + s \right) \left[ A^*D - AD^* + 2m_l (B^*C - BC^*) \right].$$

The $d\Gamma/ds$ appearing in the above equation is the one given in Eq. and $\lambda(m^2_B, m^2_{K^*_0}, s)$ is the same as that defined in Eq.

**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. 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 $(m_b/m_t)\tan^3 \beta$ ($h = h^0$, $A^0$). Here, one $\tan \beta$
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 | $C_{7}^{eff}$ | $C_{7}^{eff}$ | $C_{9}$ | $C_{9}'$ | $C_{10}$ | $C_{10}'$ |
|---------------------|--------------|--------------|--------|---------|---------|---------|
| 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       |

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 $\tau$ case.

| Wilson Coefficients | $C_{Q1}$ | $C_{Q1}'$ | $C_{Q2}$ | $C_{Q2}'$ |
|---------------------|---------|---------|---------|---------|
| 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) |

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 $C_{Q1}^{(')}$ compete with $C_{i}^{(')}$ and can overwhelm $C_{i}^{(')}$ 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 $\mu$ case and $\tau$ 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 $B \rightarrow X s \gamma$ 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 $b \rightarrow st^\pm l^-$ due to change of the sign of $C_7$ from positive to negative, while the constraint on $b \rightarrow s \gamma$ 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
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.

Higgs and charginos cancel each other. We hope that these scalar mode decays of $B$ mesons can be used to distinguish between these two regions of SUSY.

The numerical results for the decay rates, forward-backward asymmetries and polarization asymmetries of the lepton are presented in Figs. 1-5. Fig. 1 describes the differential decay rate of $B \to K^*(1430) l^+ l^-$, from 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 $C_{7}^\text{eff}$; 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. Similar effects can also be seen for the tauon case in Fig. 1b.

In Fig. 2, the forward-backward asymmetries for $B \to K^*_0 l^+ l^-$ ($l = \mu, \tau$) 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. (30) where the auxiliary function $D$ corresponds to the contributions from NHBs. Fig. 2a 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. 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.

Fig. 3(a,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 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 $C_{7}^{eфф}, C_{10}$. In large $q^2$ region, the longitudinal lepton polarization approaches to zero in all the models including the SM, because the factor $\lambda(m^2_B, m^2_{K^*}, s)$ 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 $B \rightarrow K^*_0 l^+ l^-$ on the momentum transfer squared are presented in Fig. 4. In terms of Eq. (36), one can observe that this asymmetry is sensitive to the contribution of NHBs in almost all the supersymmetric models. Fig. 4a shows the normal lepton polarization for $B \rightarrow K^*_0 \mu^+ \mu^-$. It
FIG. 4: Normal lepton polarization asymmetries for the $B \rightarrow K^*_0 l^+ l^-$ ($l = \mu, \tau$) decays as functions of $q^2$.

FIG. 5: Transverse lepton polarization asymmetries for the $B \rightarrow K^*_0 l^+ l^-$ ($l = \mu, \tau$) decays as functions of $q^2$.

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 $C^\text{eff}_7$ and $C^\text{eff}_9$. 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. 4b. Again due to same reasons as for the muons case, the normal lepton polarization changes its sign in SUSY III model.

Fig. 5 shows the dependence of transverse polarization asymmetries for $B \rightarrow K^*_0 l^+ l^-$ on the square of momentum transfer. From Eq. (37) 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 non-diagonal 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 $B \rightarrow K^*_0 l^+ l^-$ 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. 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 LHC and BTeV machines, 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.

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:

- 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.

- 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.

- 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.

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