Double-lepton polarization asymmetries in $B \to K_0^*(1430) l^+ l^-$ decay in the fourth-generation standard model

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In this paper, by deriving the expressions for double-lepton polarization asymmetries for B to a scalar meson transition in SM and SM4 and considering the corresponding uncertainties in SM we investigate the indirect effects of the fourth generation of quarks on such asymmetries in the $B \to K_0^*(1430)\ell^+\ell^-$ decay. We also compare these asymmetries with those of $B \to K\ell^+\ell^-$ decay and find out that most of these asymmetries behave similarly to the corresponding asymmetries for $B \to K\ell^+\ell^-$ decay. We finally show that in the μ channel all asymmetries, except $\langle P_{LL}\rangle$, and in τ channel only $\langle P_{LN}\rangle$ can manifest the influence of the fourth generation at the minimum value of $m_{t'}$ around 200 GeV. In addition, it is shown that for the τ channel the asymmetries such as $\langle P_{LL}\rangle$, $\langle P_{NN}\rangle$, and $\langle P_{TT}\rangle$ can indicate the effect of such new physics at $m_{t'} \geq 300$ GeV. Hence, the $B \to K_0^*(1430)\ell^+\ell^-$ decay is a valuable tool for probing new physics beyond SM, especially in the indirect searches of the fourth generation of quarks (t', b') via its manifestations in loop diagrams.

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I. INTRODUCTION

Lepton flavor violating transitions and flavor changing neutral current transitions provide an excellent testing ground for the SM. These decays, which are forbidden in the standard model (SM) at tree level, occur at loop level and are very sensitive to the gauge structure of the SM. Besides, these decays are quite sensitive to the existence of new physics beyond the SM, since loops with new particles can give considerable contribution to rare decays.

Since there exists no theoretical restrictions for the number of generations in the SM, it is possible to introduce a new generation of quarks. As a result, one possible extension of SM could be an extra fourth generation of quarks (SM4). On the other hand, the existence of a new family of quarks with the mass more than the half of the mass of the Z boson are not excluded from the CERN LEP II experiment [1] and the electroweak precision data favors an extra generation of heavy quarks, if the mass difference between the new up- and down-type quarks obeys the following relation [2]:

$$m_{t'} - m_{b'} \simeq \left(1 + \frac{1}{5}\ln\frac{m_H}{115}\right) \times 50. \tag{1}$$

In addition, the flavor democracy in the three generations of the SM [3] supports SM4. In this scenario, the masses of the first three fermion families, as well as intergenerational mixing are generated by small braking of flavor democracy [4,5]. The fourth-generation quarks possess family members with almost the same mass scale which is constrained by the experimental value of ρ and S parameters. Considering the latest data $\rho = 1.0002^{+0.0007}_{-0.0004}$ [6], the mass of the fourth quark $m_{t'}$ lies between 320 GeV and 730 GeV. If such exotic quarks exist then it could be produced by gluon fusion mechanism at LHC.

The indirect effects of the fourth-generation scenario on the rare B meson decays which include flavor changing neutral current transitions, $b \to s(d)$ transitions, have already been investigated by many researchers. The fourth generation can affect physical observables, i.e., branching ratio, CP asymmetry, polarization asymmetries, and forward-backward asymmetries. The study of these physical observables can be a good tool to look for the fourth generation of up-type quarks [7–24]. For example, it is demonstrated that the fourth-generation quarks can resolve CP violation problem in nonleptonic decays of B which are penguin-dominated and substantially reduce the differences in our findings between theory and experiment [25]. The sequential fourth generation of up quarks (t'), like u, c, t quarks, contributes to the $b \to s(d)$ transition at loop level. Hence, this new generation will change only the values of the Wilson coefficients via the virtual exchange of the fourth-generation up quark t' and the full operator set is exactly the same as in SM.

Recently, the sensitivity of the double-lepton polarization asymmetries to the fourth generation in the transition of B to a pseudo scalar meson ($B \to K\ell^+\ell^-$) [18] and B_s to a vector meson ($B_s \to \phi\ell^+\ell^-$) [19] have been investigated and it is found out that these asymmetries are very sensitive to the fourth-generation parameters ($m_{t'}$, $V_{t'b}V^*_{t's}$). In this work, by considering the theoretical and experimental uncertainties in SM we investigate the effects of the fourth-generation of quarks (b', t') on the double-lepton polarizations in the transition of B to a scalar meson ($B \to K_0^*(1430)\ell^+\ell^-$) and our results also compared to those of $B \to K\ell^+\ell^-$ decay presented in Ref. [18].

This paper is organized as follows. In Sec. II, the expressions for the matrix elements of B to a scalar meson,

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here $B \to K_0^*(1430)\ell^+\ell^-$, in SM and SM4 have been derived. In Sec. III, the general expressions for the double-lepton polarization asymmetries have been calculated. The sensitivity of these polarizations to the fourth-generation parameters $(m_{t'}, r_{sb}, \phi_{sb})$ have been numerically analyzed in Sec. IV. In the final section, a summary of conclusions is presented.

II. STRATEGY

In this section we present the theoretical expressions for the decay widths within SM4. As it has been noted above, t' like u, c, t quarks, contribute $b \to s(d)$ transition at loop level. As a result only the Wilson coefficients in SM are modified and the full operator set is exactly the same as in SM. Therefore, the relevant effective Hamiltonian for $B \to K_0^*\ell^+\ell^-$ decay, which is described by $b \to s\ell^+\ell^-$ transition at quark level, can be written as

$$\mathcal{H}_{\rm eff} = -\frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* \sum_{i=1}^{10} C_i(\mu) \mathcal{O}_i(\mu), \tag{2}$$

where the Wilson coefficients are modified as follows:

$$\lambda_t C_i \to \lambda_t C_i^{\rm SM} + \lambda_{t'} C_i^{\rm new}, \qquad \lambda_f = V_{fb} V_{fs}^*, \tag{3}$$

and the complete set of the operators $\mathcal{O}_i$ as well as the corresponding expressions for the Wilson coefficients C_i in SM are given in [26]. Also, the explicit forms of $C_i^{\rm new}$ can easily be obtained from the corresponding Wilson coefficients in SM by replacing $m_t \to m_{t'}$ (where m_t and $m_{t'}$ are the masses of quarks t and t', respectively) [26] and $\lambda_{t'}$ can be parametrized as $\lambda_{t'} = r_{sb} e^{i\phi_{sb}}$. Now, using the above effective Hamiltonian, the one-loop matrix elements of $b \to s\ell^+\ell^-$ can be given in terms of the tree-level matrix elements of the effective operators as

$$\begin{aligned}\mathcal{M}(b\to s\ell^+\ell^-) &= \langle s\ell^+\ell^-|\mathcal{H}_{\rm eff}|b\rangle \\ &= -\frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* \sum_i C_i^{\rm eff}(\mu) \langle s\ell^+\ell^-|\mathcal{O}_i|b\rangle^{\rm tree}, \\ &= -\frac{G_F \alpha}{2\pi\sqrt{2}} V_{tb} V_{ts}^* [\tilde{C}_9^{\rm eff} \bar{s}\gamma_\mu(1-\gamma_5) b \bar{\ell}\gamma_\mu \ell \\ &\quad + \tilde{C}_{10}^{\rm eff} \bar{s}\gamma_\mu(1-\gamma_5) b \bar{\ell}\gamma_\mu\gamma_5 \ell \\ &\quad - 2C_7^{\rm eff} \frac{m_b}{q^2} \bar{s} i\sigma_{\mu\nu} q^\nu (1+\gamma_5) b \bar{\ell}\gamma_\mu \ell],\end{aligned} \tag{4}$$

where $q^2 = (p_1 + p_2)^2$ and p_1 and p_2 are the final leptons four-momenta. The effective Wilson coefficients in SM are presented in [26,27] and the modified effective Wilson coefficients are as follows:

$$\begin{aligned}C_i^{\rm eff}(\mu) &= C_i^{\rm eff\ SM}(\mu) + \frac{\lambda_{t'}}{\lambda_t} C_i^{\rm eff\ new}(\mu), \qquad i = 7, \\ \tilde{C}_i^{\rm eff}(\mu) &= \tilde{C}_i^{\rm eff\ SM}(\mu) + \frac{\lambda_{t'}}{\lambda_t} \tilde{C}_i^{\rm eff\ new}(\mu), \qquad i = 9, 10.\end{aligned} \tag{5}$$

Like $C_i^{\rm new}$ the explicit forms of $C_i^{\rm eff\ new}$ and $\tilde{C}_i^{\rm eff\ new}$ can also be found from the corresponding Wilson coefficients in SM by substituting $m_t \to m_{t'}$ [26]. It should be noted that in the effective coefficient $C_9^{\rm eff}$,

$$C_9^{\rm eff} = \frac{\alpha}{2\pi} \tilde{C}_9^{\rm eff} = C_9 + \frac{\alpha}{2\pi} Y(s), \tag{6}$$

we neglect the effect of long-distance contributions coming from the real $c\bar{c}$ intermediate states for simplicity, and only consider the effect of short-distance contributions due to the one-loop matrix element of the four quark operators which is given by

$$\begin{aligned}Y_{\rm per}(s) = g\left(\frac{m_c}{m_b}, s\right)&(3C_1 + C_2 + 3C_3 + C_4 + 3C_5 + C_6) \\ &- \frac{1}{2} g(1, s)(4C_3 + 4C_4 + 3C_5 + C_6) - \frac{1}{2} g(0, s) \\ &\times (C_3 + 3C_4) + \frac{2}{9}(3C_3 + C_4 + 3C_5 + C_6),\end{aligned} \tag{7}$$

where the explicit expressions for the g functions can be found in [26].

At this stage, it is worth noting that the Glashow-Iliopoulos-Maiani mechanism dominates all the above equations. We can check the validity of this mechanism by using the unitarity property of the 4×4 Cabbibo-Kobayoshi-Maskawa matrix which leads to

$$\lambda_u + \lambda_c + \lambda_t + \lambda_{t'} = 0. \tag{8}$$

By following the steps below

$$\begin{aligned}\lambda_t C_i &= \lambda_t C_i^{\rm SM} + \lambda_{t'} C_i^{\rm new} \\ &= -(\lambda_u + \lambda_c) C_i^{\rm SM} + \lambda_{t'}(C_i^{\rm new} - C_i^{\rm SM}) \\ &= -(\lambda_u + \lambda_c) C_i^{\rm SM} \\ &= \lambda_t C_i^{\rm SM},\end{aligned} \tag{9}$$

one can finally see that the factor $\lambda_t C_i$ should be modified to $\lambda_t C_i^{\rm SM}$ when $m_{t'} \to m_t$ or $\lambda_{t'} \to 0$.

Now, having the matrix element, describing the $b \to s\ell^+\ell^-$ transition we can write down the matrix element for the $B \to K_0^*\ell^+\ell^-$ decay. In order to obtain this quantity we need to sandwich the matrix elements in Eq. (4) between the final and initial meson states. It follows from Eq. (4) that only the matrix elements $\langle K_0^*|\bar{s}\gamma_\mu(1-\gamma_5)b|B\rangle$ and $\langle K_0^*|\bar{s}i\sigma_{\mu\nu}q^\nu(1+\gamma_5)b|B\rangle$ are needed which can be parametrized in terms of form factors as follows:

$$\begin{aligned}\langle K_0^*(1430)|\bar{s}\gamma_\mu\gamma_5 b|B\rangle = f_+(q^2)&(p_B + p_{K_0^*})_\mu \\ &+ f_-(q^2) q_\mu,\end{aligned} \tag{10}$$

$$\langle K_0^*(1430)|\bar{s}i\sigma_{\mu\nu}q^\nu\gamma_5 b|B\rangle = \frac{f_T(q^2)}{m_B + m_{K_0^*}}[(p_B + p_{K_0^*})_\mu q^2 - (m_B^2 - m_{K_0^*}^2)q_\mu], \quad (11)$$

where $q_\mu = (p_B - p_{K_0^*})_\mu$.

As it can be seen from the above equations, we have to compute the form factors to evaluate the physical quantities at hadronic level. For the estimation of the form factors a nonperturbative approach is needed. In the present work, we use the three-point QCD sum rules predictions for the relevant form factors of the $B \to K_0^*$ transition in which the form factors

$$F(q^2) \in \{f_+(q^2), f_-(q^2), f_T(q^2)\},$$

are fitted to the following functions [28]:

$$F(q^2) = \frac{F(0)}{1 - a_F\frac{q^2}{m_B^2} + b_F(\frac{q^2}{m_B^2})^2}, \quad (12)$$

where the parameters $F(0)$, a_F, and b_F are listed in Table I.

Considering the above equations the transition operator is calculated as

$$\begin{aligned}\mathcal{M}(B \to K_0^*\ell^+\ell^-) &= \frac{G_F\alpha_{em}}{4\sqrt{2}\pi}V_{tb}V_{ts}^*\{\bar{\ell}\gamma^\mu\ell[A(p_B + p_{K_0^*})_\mu + Bq_\mu] \\ &\quad + \bar{\ell}\gamma^\mu\gamma_5\ell[C(p_B + p_{K_0^*})_\mu + Dq_\mu]\}. \end{aligned} \quad (13)$$

The functions appearing in Eq. (13) are defined as

$$\begin{aligned} A &= -(C_{LL} + C_{LR})f_+ + 2(C_{BR} - C_{SL})\frac{f_T}{m_B + m_{K_0^*}}, \\ B &= -(C_{LL} + C_{LR})f_- - 2(C_{BR} - C_{SL}) \\ &\quad \times \frac{f_T}{(m_B + m_{K_0^*})q^2}(m_B^2 - m_{K_0^*}^2), \\ C &= -(C_{LR} - C_{LL})f_+, \\ D &= -(C_{LR} - C_{LL})f_-, \end{aligned}$$

where

$$\begin{aligned} C_{LL} &= (\tilde{C}_9^{\rm eff} - \tilde{C}_{10}^{\rm eff}), \qquad C_{LR} = (\tilde{C}_9^{\rm eff} + \tilde{C}_{10}^{\rm eff}), \\ C_{SL} &= -2m_s C_7^{\rm eff}, \qquad C_{BR} = -2m_b C_7^{\rm eff}. \end{aligned} \quad (14)$$

TABLE I. The form factors for $B \to K_0^*(1430)\ell^+\ell^-$ in a three-parameter fit [28].

	$F(0)$	a_F	b_F
$f_+^{B\to K_0^*}$	0.31 ± 0.08	0.81	-0.21
$f_-^{B\to K_0^*}$	-0.31 ± 0.07	0.80	-0.36
$f_T^{B\to K_0^*}$	-0.26 ± 0.07	0.41	-0.32

From the above equations we get the following result for the differential decay width:

$$\frac{d\Gamma}{d\hat{s}}(B \to K_0^*\ell^+\ell^-) = \frac{G^2\alpha^2 m_B}{2^{14}\pi^5}|V_{tb}V_{ts}^*|^2\lambda^{1/2}(1, \hat{r}_{K_0^*}, \hat{s})v\Delta(\hat{s}), \quad (15)$$

with

$$\begin{aligned}\Delta(\hat{s}) &= \frac{4m_B^2}{3}\,\mathrm{Re}[24m_B^2\hat{m}_l^2(1 - \hat{r}_{K_0^*})D^\star C \\ &\quad + \lambda m_B^2(3 - v^2)|A|^2 + 12m_B^2\hat{m}_l^2\hat{s}|D|^2 \\ &\quad + m_B^2|C|^2\{2\lambda - (1 - v^2)(2\lambda - 3(1 - \hat{r}_{K_0^*})^2)\}], \end{aligned} \quad (16)$$

where $\hat{s} = q^2/m_B^2$, $\hat{r}_{K_0^*} = m_{K_0^*}^2/m_B^2$ and $\lambda(a, b, c) = a^2 + b^2 + c^2 - 2ab - 2ac - 2bc$, $\hat{m}_\ell = m_\ell/m_B$, $v = \sqrt{1 - 4\hat{m}_\ell^2/\hat{s}}$ is the final lepton velocity.

III. DOUBLE-LEPTON POLARIZATION

Having obtained the matrix elements for the $B \to K_0^*\ell^+\ell^-$ decay in the last section, we can now calculate the double-polarization asymmetries. For this, we define the orthogonal unit vectors $s_i^{\pm\mu}$ in the rest frame of leptons as the following:

$$\begin{aligned} s_L^{-\mu} &= (0, \vec{e}_L^{\,-}) = \left(0, \frac{\vec{p}_-}{|\vec{p}_-|}\right), \\ s_L^{+\mu} &= (0, \vec{e}_L^{\,+}) = \left(0, \frac{\vec{p}_+}{|\vec{p}_+|}\right), \\ s_N^{-\mu} &= (0, \vec{e}_N^{\,-}) = \left(0, \frac{\vec{p}_{K_0^*} \times \vec{p}_-}{|\vec{p}_{K_0^*} \times \vec{p}_-|}\right), \\ s_N^{+\mu} &= (0, \vec{e}_N^{\,+}) = \left(0, \frac{\vec{p}_{K_0^*} \times \vec{p}_+}{|\vec{p}_{K_0^*} \times \vec{p}_+|}\right), \\ s_T^{-\mu} &= (0, \vec{e}_T^{\,-}) = (0, \vec{e}_N^{\,-} \times \vec{e}_L^{\,-}), \\ s_T^{+\mu} &= (0, \vec{e}_T^{\,+}) = (0, \vec{e}_N^{\,+} \times \vec{e}_L^{\,+}), \end{aligned} \quad (17)$$

where $i = L, N$ and T are the aberrations of longitudinal, normal and transversal polarization directions and $\vec{p}_\mp$ and $\vec{p}_{K_0^*}$ are the three—momenta of the leptons $\ell^\mp$ and K_0^* meson, respectively. Now by using Lorentz transformation to boost the unit vectors from the rest frame of the leptons to the center of mass (CM) frame of leptons one finds that only the longitudinal unit vectors $s_L^{\pm\mu}$ change as shown below:

$$(s_L^{-\mu})_{\rm CM} = \left(\frac{|\vec{p}_-|}{m_\ell}, \frac{E\vec{p}_-}{m_\ell|\vec{p}_-|}\right),$$
$$(s_L^{+\mu})_{\rm CM} = \left(\frac{|\vec{p}_-|}{m_\ell}, -\frac{E\vec{p}_-}{m_\ell|\vec{p}_-|}\right), \tag{18}$$

and the other two cases remain unchanged. The polarization asymmetries can now be calculated using the spin projector $\frac{1}{2}(1+\gamma_5 \not{s}_i^-)$ for ℓ^- and the spin projector $\frac{1}{2}(1+\gamma_5 \not{s}_i^+)$ for ℓ^+.

Considering the above explanations, we can define the double-lepton polarization asymmetries as in [29]:

$$P_{ij}(\hat{s}) = \frac{\left(\frac{d\Gamma}{d\hat{s}}(\vec{s}_i^-, \vec{s}_j^+) - \frac{d\Gamma}{d\hat{s}}(-\vec{s}_i^-, \vec{s}_j^+)\right) - \left(\frac{d\Gamma}{d\hat{s}}(\vec{s}_i^-, -\vec{s}_j^+) - \frac{d\Gamma}{d\hat{s}}(-\vec{s}_i^-, -\vec{s}_j^+)\right)}{\left(\frac{d\Gamma}{d\hat{s}}(\vec{s}_i^-, \vec{s}_j^+) + \frac{d\Gamma}{d\hat{s}}(-\vec{s}_i^-, \vec{s}_j^+)\right) + \left(\frac{d\Gamma}{d\hat{s}}(\vec{s}_i^-, -\vec{s}_j^+) + \frac{d\Gamma}{d\hat{s}}(-\vec{s}_i^-, -\vec{s}_j^+)\right)}, \tag{19}$$

where $i, j = L, N, T$, and the first index i corresponds to the lepton while the second index j corresponds to the antilepton, respectively. After doing the straightforward calculations we obtain the following expressions for $P_{ij}(\hat{s})$:

$$P_{LL} = \frac{-4m_B^2}{3\Delta}\,\mathrm{Re}[-24m_B^2\hat{m}_l^2(1-\hat{r}_{K_0^*})C^\star D + \lambda m_B^2(1+v^2)|A|^2 \tag{20}$$

$$-12m_B^2\hat{m}_l^2\hat{s}|D|^2 + m_B^2|C|^2(2\lambda-(1-v^2)\times(2\lambda+3(1-\hat{r}_{K_0^*})^2))], \tag{21}$$

$$P_{LN} = \frac{-4\pi m_B^3\sqrt{\lambda\hat{s}}}{\hat{s}\Delta}\,\mathrm{Im}[-m_B\hat{m}_l\hat{s}A^\star D - m_B\hat{m}_l(1-\hat{r}_{K_0^*})A^\star C], \tag{22}$$

$$P_{NL} = -P_{LN}, \tag{23}$$

$$P_{LT} = \frac{4\pi m_B^3\sqrt{\lambda\hat{s}}}{\hat{s}\Delta}\,\mathrm{Re}[m_B\hat{m}_l v(1-\hat{r}_{K_0^*})|C|^2 + m_B\hat{m}_l v\hat{s}C^\star D], \tag{24}$$

$$P_{TL} = P_{LT}, \tag{25}$$

$$P_{NT} = -\frac{8m_B^2 v}{3\Delta}\,\mathrm{Im}[2\lambda m_B^2 A^\star C], \tag{26}$$

$$P_{TN} = -P_{NT}, \tag{27}$$

$$P_{TT} = \frac{4m_B^2}{3\Delta}\,\mathrm{Re}[-24m_B^2\hat{m}_l^2(1-\hat{r}_{K_0^*})C^\star D - \lambda m_B^2 \times(1+v^2)|A|^2 - 12m_B^2\hat{m}_l^2\hat{s}|D|^2 + m_B^2|C|^2 \times\{2\lambda-(1-v^2)(2\lambda+3(1-\hat{r}_{K_0^*})^2)\}], \tag{28}$$

$$P_{NN} = \frac{4m_B^2}{3\Delta}\,\mathrm{Re}[24m_B^2\hat{m}_l^2(1-\hat{r}_{K_0^*})C^\star D - \lambda m_B^2(3-v^2) \times|A|^2 + 12m_B^2\hat{m}_l^2\hat{s}|D|^2 + m_B^2|C|^2\{2\lambda-(1-v^2) \times(2\lambda-3(1-\hat{r}_{K_0^*})^2)\}]. \tag{29}$$

IV. RESULTS AND DISCUSSIONS

In this section the analytical dependence of the double-lepton polarizations on the fourth quark mass ($m_{t'}$) and the product of quark mixing matrix elements ($V_{t'b}^* V_{t's} = r_{sb}e^{i\phi_{sb}}$) are studied. As pointed out in Sec. II, for the main input parameters which are form factors we have chosen the predictions of the three-point QCD sum rules method [28]. For the other input parameters we use the following values [30]:

$$\begin{aligned}
&m_B = 5.279 \pm 0.03\ \mathrm{GeV}, \qquad m_{K_0^*} = 1.425 \pm 0.05\ \mathrm{GeV},\\
&m_b = 4.19^{+0.18}_{-0.06}\ \mathrm{GeV}, \qquad m_c = 1.27^{+0.07}_{-0.09}\ \mathrm{GeV},\\
&m_s = 0.101^{+0.029}_{-0.021}\ \mathrm{GeV}, \qquad m_\mu = 0.105\ \mathrm{GeV},\\
&m_\tau = 1.77\ \mathrm{GeV}, \qquad \alpha^{-1} = 129,\\
&\tau_B = (1.525 \pm 0.009)\times 10^{-12}\ \mathrm{s},\\
&\lambda = 0.2253 \pm 0.0007, \qquad A = 0.808^{+0.022}_{-0.015},\\
&\bar{\rho} = 0.132^{+0.022}_{-0.014}, \qquad \bar{\eta} = 0.341 \pm 0.013,
\end{aligned} \tag{30}$$

where A, λ, $\bar{\rho}$, $\bar{\eta}$ are the Wolfenstein parameters in the Cabbibo-Kobayoshi-Maskawa matrix. In order to present a quantitative analysis of the double-lepton polarization asymmetries, the values of fourth-generation parameters are needed. Considering the experimental values of $B \to X_s\gamma$ and $B \to X_s\ell^+\ell^-$ decays the value of the parameter r_{sb}, lies in the range {0.01–0.03} for $\phi_{sb} \sim \{0°–360°\}$ and $m_{t'} \sim \{200–600\}$ GeV [11,22]. Using the B_s mixing parameter Δm_{B_s}, a sharp limit on ϕ_{sb} has been obtained around 90° [7]. Therefore in our upcoming numerical analysis, the corresponding values of above ranges are: $r_{sb} = \{0.01, 0.02, 0.03\}$, $\phi_{sb} = \{60°, 90°, 120°\}$, $m_{t'} = 175\ \mathrm{GeV} \le m_{t'} \le 600\ \mathrm{GeV}$.

It is clear from the expressions of all nine double-lepton polarization asymmetries that they depend on the momentum transfer q^2 and the new parameters ($m_{t'}, r_{sb}, \phi_{sb}$). Consequently, it may be experimentally difficult to investigate these dependencies at the same time. In the present work, we omit the q^2 dependency by integrating over this parameter and investigating the averaged double-lepton polarization asymmetries. The average of P_{ij} over q^2 is defined as

$$\langle P_{ij}\rangle = \frac{\int_{4\hat{m}_\ell^2}^{(1-\sqrt{\hat{r}_{K_0^*}})^2} P_{ij}\frac{d\mathcal{B}}{d\hat{s}}d\hat{s}}{\int_{4\hat{m}_\ell^2}^{(1-\sqrt{\hat{r}_{K_0^*}})^2} \frac{d\mathcal{B}}{d\hat{s}}d\hat{s}}. \tag{31}$$

Using the above formula we have presented our analysis for the dependency of $\langle P_{ij}\rangle$ on the fourth-generation parameters in a series of figures (see Figs. 1–6). As it is seen from these figures, the SM4 diagrams of each asymmetry cut the corresponding SM diagram at $m_{t'} = m_t$ which is consistent with the result of the Glashow-Iliopoulos-Maiani mechanism appearing in Eq. (9). By considering the theoretical and experimental uncertainties for $B \to K_0^* \ell^+\ell^-$ decay in SM, we compare our results for $B \to K_0^* \ell^+\ell^-$ decay with the results of Ref. [18] for $B \to K\ell^+\ell^-$ decay. It should be mentioned that the theoretical uncertainties come from the hadronic uncertainties related to the form factors and the experimental uncertainties originate from the mass of quarks and hadrons and Wolfestein parameters. Our study shows that the overall behavior of $\langle P_{ij}\rangle$ with regard to $m_{t'}$, r_{sb}, and ϕ_{sb} are, to a large extent, the same as that of $B \to K\ell^+\ell^-$ decay, therefore, we only discuss the differences of these two decay modes and investigate some related topics not discussed in Ref. [18]:

(i) *Figure* 1: Similar to the $B \to K$ decay, the maximum deviation of $\langle P_{LL}\rangle$ for the μ channel from the SM value is less than 5%, therefore we do not present the relevant figures for this channel. However, as seen from the plots in Fig. 1 for the τ channel of this decay, the maximum deviation of such quantity from SM is about 30% which is comparable to this value for $B \to K$ decay which is 20% [18]. Although the magnitude of this quantity does not show strong dependency to the fourth-generation parameters, it is found out from Tables II and III that the upper limit of the uncertainty of that value in the SM can be covered by its corresponding deviation from the SM

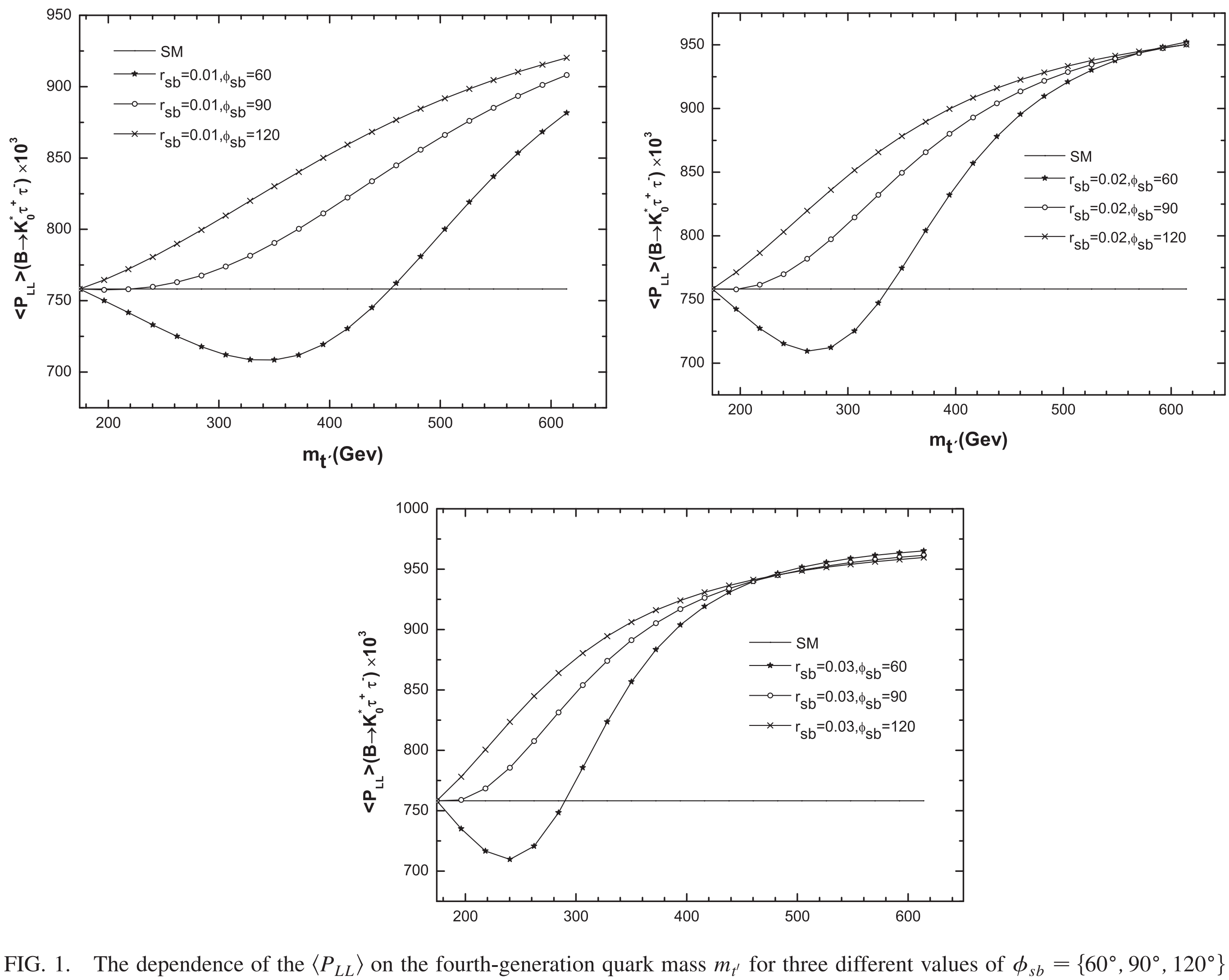

FIG. 1. The dependence of the $\langle P_{LL}\rangle$ on the fourth-generation quark mass $m_{t'}$ for three different values of $\phi_{sb} = \{60°, 90°, 120°\}$ and $r_{sb} = \{0.01, 0.02, 0.03\}$ for the τ channel.

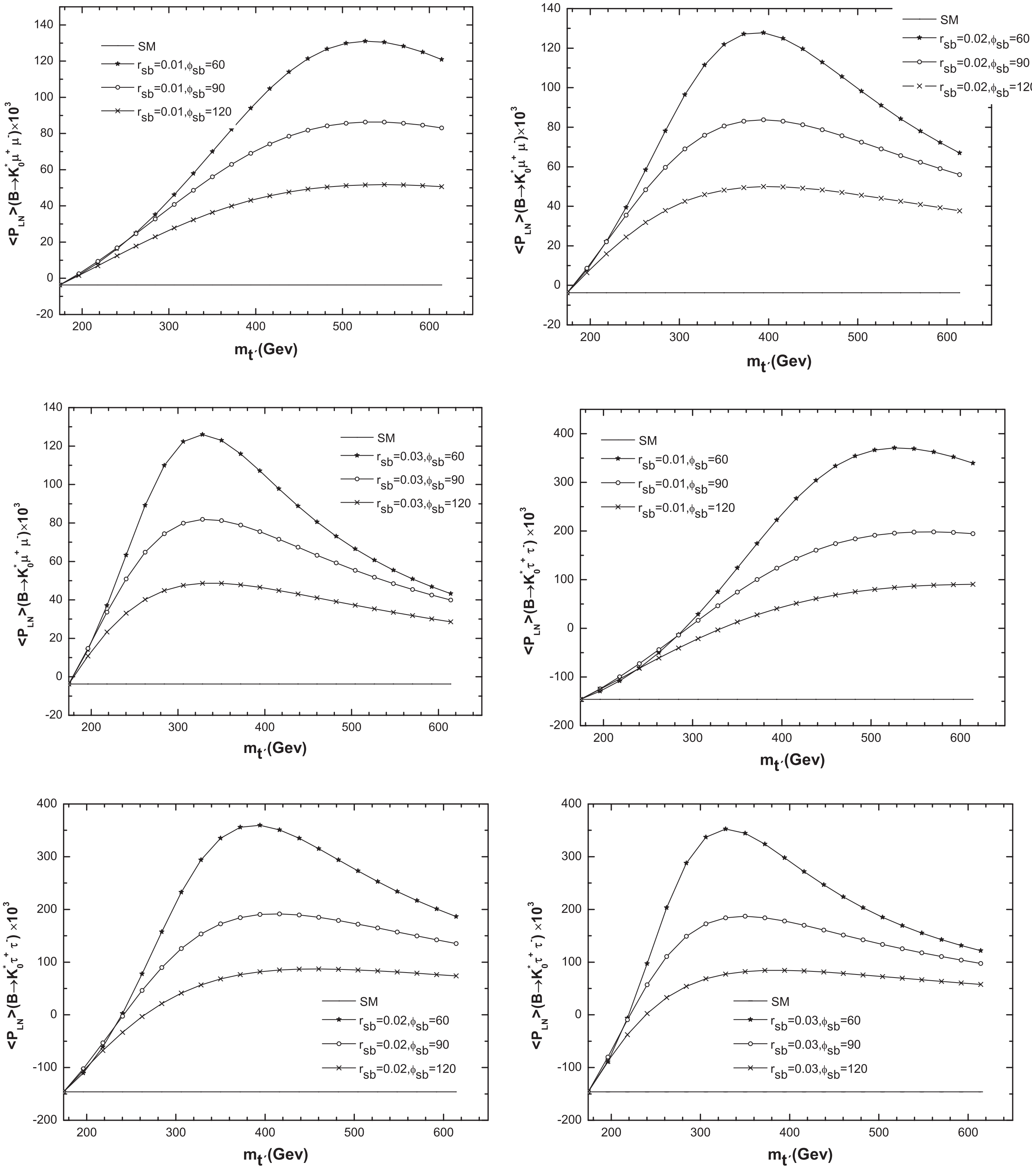

FIG. 2. The dependence of the $\langle P_{LN}\rangle$ on the fourth-generation quark mass $m_{t'}$ for three different values of $\phi_{sb} = \{60°, 90°, 120°\}$ and $r_{sb} = \{0.01, 0.02, 0.03\}$ for the μ and τ channels.

in some parameter ranges. For example, it is understood from the three plots in Fig. 1 that by increasing r_{sb} and ϕ_{sb}, simultaneously or not, the lower limit of $m_{t'}$, where the deviation from the SM exceeds from the corresponding SM uncertainty, decreases which happens at $m_{t'} \approx 300$ Gev for $r_{sb} = 0.03$ and $\phi_{sb} = 120°$.

(ii) *Figure* 2: From these plots, it is apparent that the magnitude of $\langle P_{LN}\rangle$ in $B \to K_0^*$ exhibits strong dependence on the fourth-generation parameters,

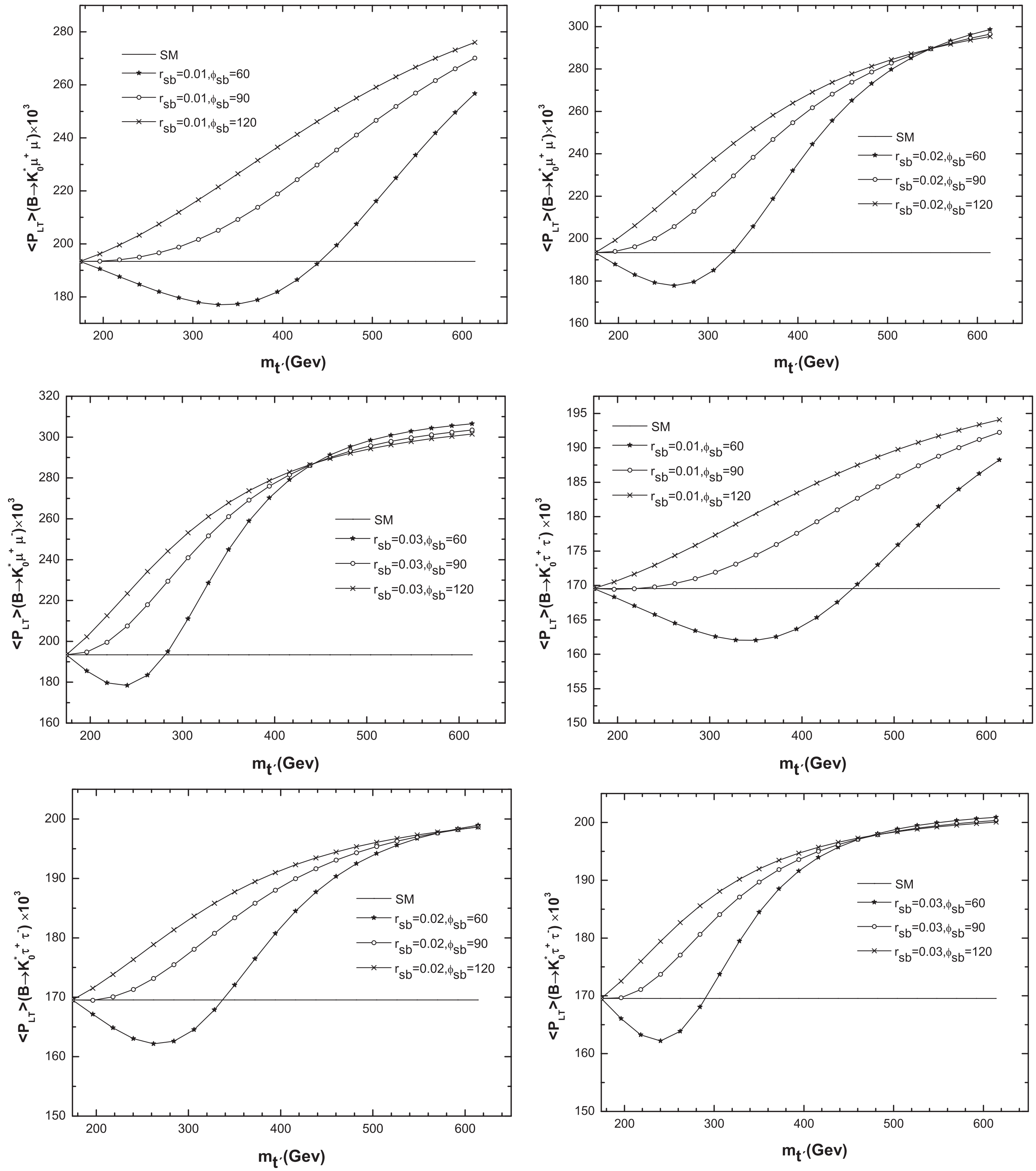

FIG. 3. The dependence of the $\langle P_{LT}\rangle$ on the fourth-generation quark mass $m_{t'}$ for three different values of $\phi_{sb} = \{60°, 90°, 120°\}$ and $r_{sb} = \{0.01, 0.02, 0.03\}$ for the μ and τ channels.

such that it becomes approximately 2000 times greater than the SM value for the μ channel and 3 times larger than the SM value for the τ channel. It is seen from Ref. [18] that the corresponding plots for the μ and the τ channels of this and $B \to K$ processes are very similar to each other, indicating particularly that the maximum deviation for the $B \to K_0^*$ transition and the $B \to K$ transition are very close

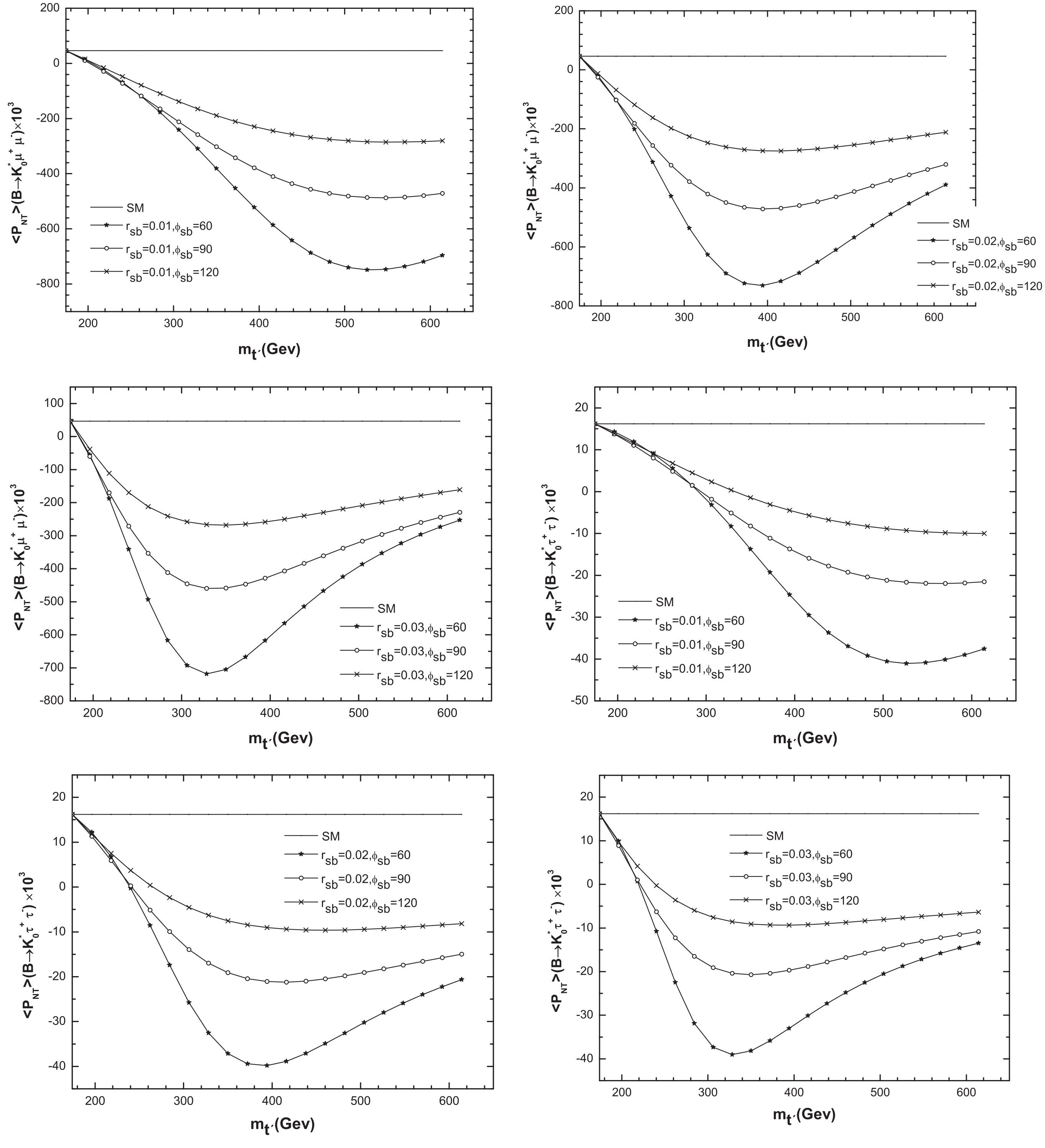

FIG. 4. The dependence of the $\langle P_{NT}\rangle$ on the fourth-generation quark mass $m_{t'}$ for three different values of $\phi_{sb} = \{60°, 90°, 120°\}$ and $r_{sb} = \{0.01, 0.02, 0.03\}$ for the μ and τ channels.

to each other. Also, from the same comparison we conclude that by considering the short-distance effects in $C_9^{\rm eff}$, coming from the one-loop matrix element of the four quark operators, the sign of $\langle P_{LN}\rangle$ in SM4 may change compared to the SM prediction. (In Ref. [18], the short-distance contributions in $C_9^{\rm eff}$ are not considered.) On the other hand, it is seen from Tables II and III that the SM value at most reaches up to zero for the μ channel and becomes negative for the τ channel. As a result, the SM values can not interfere with the positive predictions of SM4 for $\langle P_{LN}\rangle$. Furthermore, it is understood from the plots

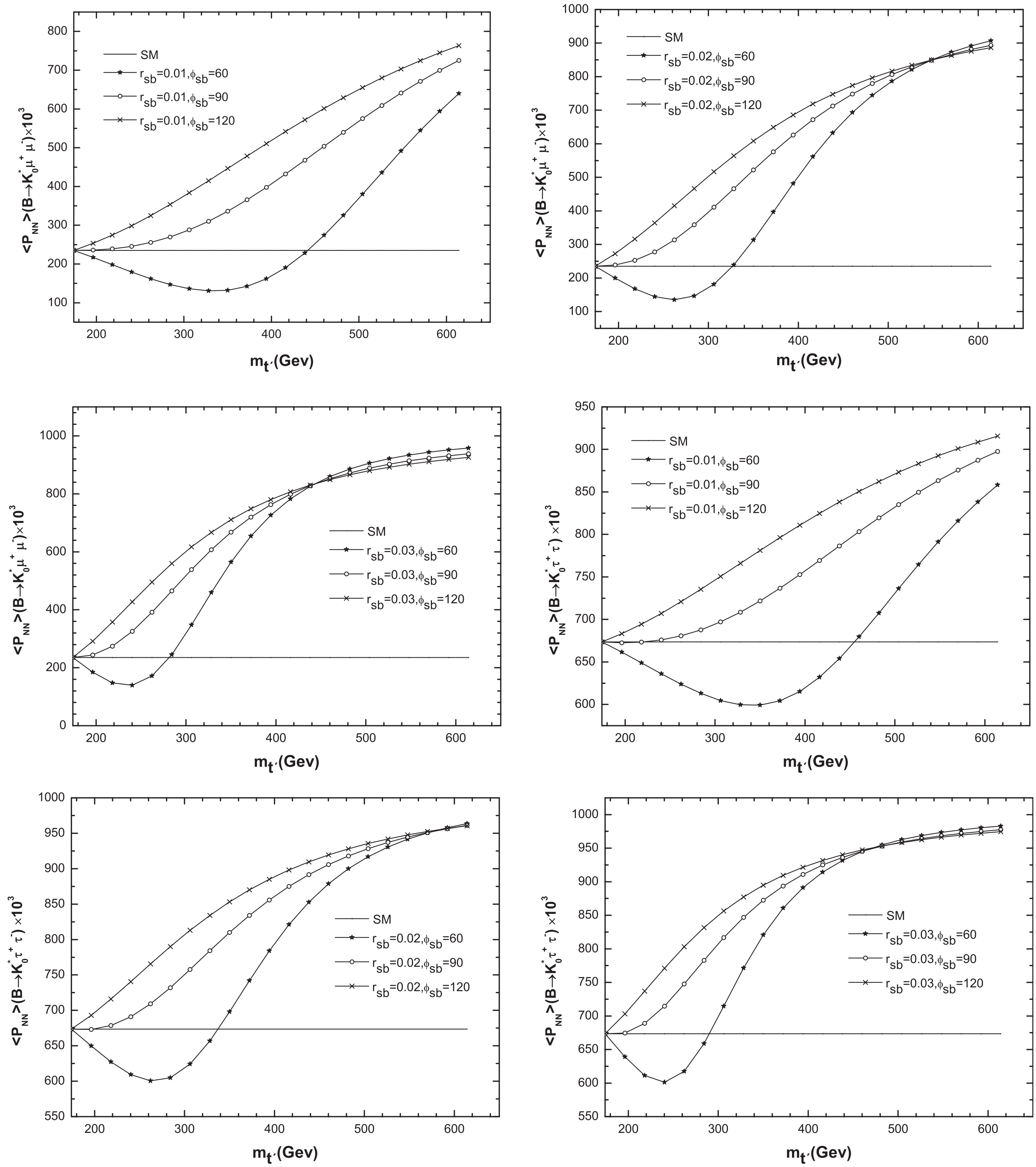

FIG. 5. The dependence of the $\langle P_{NN} \rangle$ on the fourth-generation quark mass $m_{t'}$ for three different values of $\phi_{sb} = \{60°, 90°, 120°\}$ and $r_{sb} = \{0.01, 0.02, 0.03\}$ for the μ and τ channels.

in Fig. 2 that by enhancing r_{sb}, the lower bound of $m_{t'}$, where the error of SM becomes less than the corresponding deviation from the SM, reduces happening at $180\ \text{GeV} \leq m_{t'} \leq 200\ \text{GeV}$ for the μ case and $m_{t'} \approx 200$ GeV for the τ case occurring at $r_{sb} = 0.03$.

(iii) *Figure* 3: While the magnitude of $\langle P_{LT} \rangle$ for the μ channel of $B \to K_0^*$ decay and $B \to K$ decay

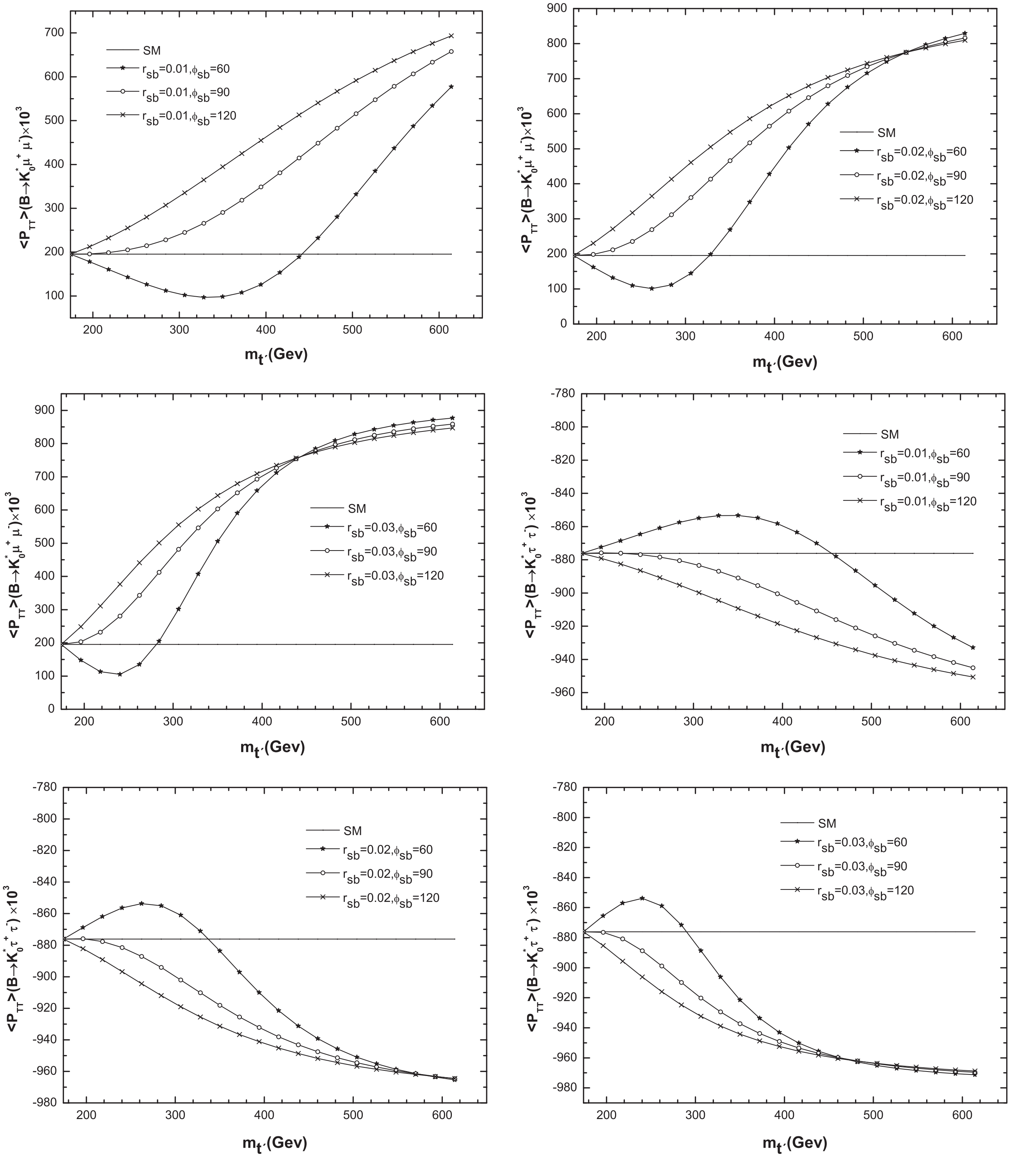

FIG. 6. The dependence of the $\langle P_{TT} \rangle$ on the fourth-generation quark mass $m_{t'}$ for three different values of $\phi_{sb} = \{60°, 90°, 120°\}$ and $r_{sb} = \{0.01, 0.02, 0.03\}$ for the μ and τ channels.

changes at most about 60% compared to the SM prediction, its value reaches up to 18% SM and 11% SM for the τ channel of $B \to K_0^*$ and $B \to K$ decay, respectively. In addition, the comparison of data in Tables II and III with the points of corresponding plots in Fig. 3 shows that although for the μ channel the upper limit of the uncertainty in the SM can be hidden by its corresponding deviation from the SM for some specific parameter ranges, for the τ channel this will not happen. As a result, for establishing the fourth generation of quarks the measurement of $\langle P_{LT}\rangle$ for the τ channel of the $B \to K_0^*$ transition is not suitable. Our analysis regarding the μ channel also implies this phenomena that by increasing r_{sb} and ϕ_{sb} the lower limit of an admissible $m_{t'}$ for discovering new physics starts to lower such that it occurs in $m_{t'} \approx 220$ GeV for $r_{sb} = 0.03$ and $\phi_{sb} = 120°$.

(iv) *Figure* 4: By comparing this figure with Fig. 2 and using the corresponding data given in Tables II and III, one can find out that the overall behavior of $\langle P_{NT}\rangle$ and $\langle P_{LN}\rangle$ are to a large extent the same. However, there are also some small differences. For example, it is obvious from Tables II and III that the lower bound of the SM value gets positive for the μ channel and becomes at least zero for the τ channel. As a result, the SM values can not overlap with the negative predictions of SM4 for $\langle P_{NT}\rangle$. As another example, a comparison between plots in Figs. 2 and 4 shows that even though the lower bound of an acceptable $m_{t'}$ for the μ case lies at the same range for both asymmetries, that is, 180 GeV $\leq m_{t'} \leq$ 200 GeV, it gets to $m_{t'} \approx 220$ GeV for the τ case for $r_{sb} = 0.03$ in Fig. 4. On the other hand, a comparison between these plots and those of $B \to K$ decay indicates that the μ channel and the τ channel of this process behaves like the τ channel and the μ channel of the B to K process, respectively. Therefore, the μ channel of B to K_0^*, which is approximately two thousand times greater than the SM value, is more sensitive to the fourth-generation parameters than its corresponding τ channel which is 2 times larger than the SM value.

(v) *Figure* 5: For the μ channel, both $B \to K_0^*$ and $B \to K$ decays show the same and considerable dependence on the fourth-generation parameters such that the constructive contribution of new physics enhances the SM magnitude of $\langle P_{NN}\rangle$ by a factor of 4. However, for the τ channel, such new physics can yield only about 50% enhancement to this asymmetry for the $B \to K_0^*$ process and 30% enhancement to the same asymmetry for the $B \to K$ transition. Furthermore, similar to the previous asymmetries, by increasing r_{sb} and ϕ_{sb} the lower limit of an admissible $m_{t'}$ decreases, happening at $m_{t'} \approx$ 200 GeV for the μ case, and $m_{t'} \approx 300$ GeV for the τ case, occurring at $r_{sb} = 0.03$ and $\phi_{sb} = 120°$.

(vi) *Figure* 6: For the μ channel, both $B \to K_0^*$ and $B \to K$ decays show similar considerable dependence on the fourth-generation parameters in such a way that the magnitude of $\langle P_{TT}\rangle$ becomes approximately 4 times greater than the SM value. Compared to the μ channel the τ channel is such that the destructive contribution of this new physics is more significant than the constructive effect which brings down the value of $\langle P_{TT}\rangle$ to 11% SM value for the $B \to K_0^*$ transition and 6% SM value for the $B \to K$ process. Such values for the τ channel indicate that the sensitivity of this asymmetry to the fourth-generation parameters is not considerable, but as can be seen from Table III the lower bound of the corresponding data in the SM can be covered by the deviation from the SM caused by this new physics in some parameter ranges. In addition, like the

TABLE II. The averaged double-lepton polarization asymmetries for $B \to K_0^*(1430)\mu^+\mu^-$ in the SM. The errors shown for each asymmetry are due to the theoretical and experimental uncertainties. The former is related to the theoretical uncertainty and the latter is due to experimental uncertainty.

$\langle P_{LL}\rangle$	$\langle P_{LN}\rangle$	$\langle P_{LT}\rangle$	$\langle P_{NT}\rangle$
$-0.937^{+0.005+0.001}_{-0.007-0.001}$	$-0.0037^{+0.0022+0.0014}_{-0.0014-0.0021}$	$0.193^{+0.020+0.005}_{-0.032-0.004}$	$0.046^{+0.001+0.021}_{-0.002-0.014}$
$\langle P_{NN}\rangle$	$\langle P_{TT}\rangle$		
$0.235^{+0.040+0.031}_{-0.065-0.028}$	$0.195^{+0.036+0.030}_{-0.060-0.026}$		

TABLE III. Same as Table II except for $B \to K_0^*(1430)\tau^+\tau^-$.

$\langle P_{LL}\rangle$	$\langle P_{LN}\rangle$	$\langle P_{LT}\rangle$	$\langle P_{NT}\rangle$
$0.758^{+0.129+0.013}_{-0.975-0.013}$	$-0.146^{+0.039+0.006}_{-0.052-0.006}$	$0.169^{+0.052+0.002}_{-0.050-0.002}$	$0.016^{+0.064+0.001}_{-0.008-0.001}$
$\langle P_{NN}\rangle$	$\langle P_{TT}\rangle$		
$0.673^{+0.174+0.020}_{-1.329-0.020}$	$-0.876^{+0.495+0.006}_{-0.066-0.006}$		

previous cases, our study shows that by increasing r_{sb} and ϕ_{sb} the lower limit of an admissible $m_{t'}$ reduces, getting to a minimum at $m_{t'} \approx 200$ GeV for the μ channel, and $m_{t'} \approx 360$ GeV for the τ channel, happening at $r_{sb} = 0.03$ and $\phi_{sb} = 120°$.

Finally, let us discuss briefly whether the lepton polarization asymmetries are measurable in experiments or not. Experimentally, to measure an asymmetry $\langle P_{ij}\rangle$ of the decay with branching ratio $\mathcal{B}$ at $n\sigma$ level, the required number of events (i.e., the number of $B\bar{B}$) is given by the formula

$$N = \frac{n^2}{\mathcal{B} s_1 s_2 \langle P_{ij}\rangle^2},$$

where s_1 and s_2 are the efficiencies of the leptons. The values of the efficiencies of the τ leptons differ from 50% to 90% for their various decay modes [31] and the error in τ-lepton polarization is approximately (10–15)% [32]. So, the error in measurement of the τ-lepton asymmetries is estimated to be about (20–30)%, and the error in obtaining the number of events is about 50%.

Based on the above expression for N, in order to detect the lepton polarization asymmetries in the μ and τ channels at 3σ level, the minimum number of required events are given by (the efficiency of the τ lepton is considered 0.5):

(i) for $B \to K_0^* \mu^+\mu^-$ decay

$$N \sim \begin{cases} 10^7 & (\text{for } \langle P_{LL}\rangle, \langle P_{NN}\rangle, \langle P_{TT}\rangle), \\ 10^8 & (\text{for } \langle P_{LT}\rangle, \langle P_{TL}\rangle, \langle P_{LN}\rangle, \langle P_{NL}\rangle), \\ 10^{10} & (\text{for } \langle P_{NT}\rangle, \langle P_{TN}\rangle), \end{cases}$$

(ii) for $B \to K_0^* \tau^+\tau^-$ decay

$$N \sim \begin{cases} 10^{10} & (\text{for } \langle P_{LL}\rangle, \langle P_{LN}\rangle, \langle P_{NL}\rangle, \langle P_{NN}\rangle, \langle P_{TT}\rangle), \\ 10^{11} & (\text{for } \langle P_{LT}\rangle, \langle P_{TL}\rangle), \\ 10^{12} & (\text{for } \langle P_{NT}\rangle, \langle P_{TN}\rangle). \end{cases}$$

Comparing the above values of N to the number of $B\bar{B}$ pairs ($\sim 10^{12}$ per year) which will be produced at LHC experiments (ATLAS, CMS, LHCb), it seems that for the μ channel all double-lepton polarizations and for the τ channel, probably $\langle P_{LL}\rangle$, $\langle P_{LN}\rangle$, $\langle P_{NL}\rangle$, $\langle P_{NN}\rangle$, and $\langle P_{TT}\rangle$ can be detected at the LHC. However, it should be mentioned that the muon polarization can be measured if we stop muon. It is very difficult to stop and measure the polarization of muon in the recent experiments. The tau polarization can be studied by looking at the decay products of tau. The measurement of tau polarization in this sense is easier than the polarization of muon.

V. SUMMARY

To sum up, in this paper by considering the theoretical and experimental uncertainties in SM we have presented an analysis of the double-lepton polarization asymmetries for $B \to K_0^* \ell^+\ell^-$ decay using the SM4 model. At the same time, we have compared our results to those of $B \to K\ell^+\ell^-$ decay and the following conclusions are obtained:

(i) Except for $\langle P_{NT}\rangle$ in both μ and τ channels, the behavior of all other asymmetries in $B \to K_0^* \ell^+\ell^-$ decay versus $m_{t'}$, r_{sb}, and ϕ_{sb} are, to a large extent, the same as those for $B \to K\ell^+\ell^-$ decay.

(ii) In the μ channel, except for $\langle P_{LL}\rangle$, which does not show any sensitivity to the fourth-generation parameters, all other asymmetries have a chance to show a sign of new physics at the minimum value of $m_{t'}$ around 200 GeV.

(iii) In the τ channel, only for the asymmetries $\langle P_{LN}\rangle$ and $\langle P_{NT}\rangle$ is there a possibility to see the effect of the fourth generation at the minimum value of $m_{t'}$, about 200 GeV. However, due to the constraints on the number of $B\bar{B}$ pairs produced at the LHC, only the $\langle P_{LN}\rangle$ value can be helpful for discovering the new generation of quarks. The asymmetries $\langle P_{LL}\rangle$, $\langle P_{NN}\rangle$, and $\langle P_{TT}\rangle$ can also be significant for finding such new physics at $m_{t'} \geq 300$ GeV.

Based on the above discussion, we found out that within some specific parameter ranges most of these asymmetries can show considerable dependency on the fourth-generation parameters which can be detected at the LHC. Therefore, considering $B \to K_0^* \ell^+\ell^-$ decay, it provides an opportunity to investigate the validity of the fourth-generation of quarks theory in near future.

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