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Effects of scalar leptoquark on semileptonic $\Lambda_b$ decays

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Abstract

We study the scalar leptoquark effects on the rare semileptonic decays of $\Lambda_b$ baryon, governed by the quark level transition $b \to s\ell^+\ell^-$. We estimate the branching ratios, forward–backward asymmetries, lepton polarisation parameters and the lepton flavour non-universality effects in these decay channels. We find significant deviations from the corresponding standard model predictions in some of the observables due to leptoquark effects. We also investigate the lepton flavour violating decays $\Lambda_b \to \Lambda L\ell^+\ell^-$, the branching ratios of which are found to be $O(10^{-10} - 10^{-9})$.

1. Introduction

The study of the rare $B$ meson decays involving flavour changing neutral current (FCNC) transitions is very crucial, as they provide sensitive probe to look for new physics (NP) beyond the standard model (SM). These decays are highly suppressed in the SM due to Glashow–Iliopoulos–Maiani mechanism and occur only through one-loop level penguin and box diagrams. Recently, several anomalies have been observed in the rare semileptonic $B$ decays mediated through the FCNC $b \to s$ transitions. The most prominent ones are the observation of $3.7\sigma$ deviation in the angular observable $P_3^+$ of $B \to K^*\mu^+\mu^-$ mode and the violation of lepton universality in the $B \to K\ell^+\ell^-$ decays at the level of $2.6\sigma$ by the LHCb experiment. In addition, LHCb has also observed significant discrepancy in the decay rates of the $B \to K^*\ell^+\ell^-$ processes [$5, 6$]. Also the decay rate of the $B \to \phi\mu^+\mu^-$ process [$7$] has $3.3\sigma$ deviation form its SM value in the low $q^2$ region. Furthermore, the observed discrepancy in the ratio of branching fractions of exclusive $B \to K(1^+)^+\ell^-$ decay and the inclusive decays into dimuon over dielectron in the full $q^2$ range [$8$] provide strong evidence of the presence of lepton non-universality.

The anomalies observed in $B \to s\ell^+\ell^-$ processes at LHCb [$1, 2, 4–7$] have attracted a lot of attention in recent times. The implications of these observations have been extensively studied both in the context of various NP models and in model independent ways [$9–13$]. These deviations which are at the level of $(2–3)\sigma$ are not statistically significant enough to provide an unambiguous signal of NP. On the other hand they are also not small enough to be ignored completely and need to be scrutinized meticulously as many different ways as possible. If indeed they really evince the smoking gun signal of some kind of NP, such effects must also show up in other decay channels involving $b \to s$ transitions, such as the corresponding $\Lambda_b$ transitions. Therefore, the study of the rare $\Lambda_b$ decays is of utmost importance to obtain an unambiguous signal of NP. Including the baryonic decay mode $\Lambda_b \to \Lambda(\to p\pi^-)\mu^+\mu^-$ in the Bayesian analysis of $|\Delta B| = |\Delta S| = 1$ transitions, a fit of the Wilson coefficients $C_{9,10}, C'_{9,10}$ has been performed in [$14$], and it has been shown that, the shift to $C_9$ prefers to be opposite to the one found in mesonic case. To be more specific, the shift in $C_9$ in baryonic decay is found to be $\Delta_9 = C_9 - C_9^{\text{SM}} = 1.6_{-0.9}^{+0.7}$, as compared to the mesonic case where its value is $\Delta_9 = -1.09_{-0.20}^{+0.32}$ [$11$]. Whereas the corresponding shifts in $C_{10}$ are in the same direction, i.e., $\Delta_{10} = 0.7_{-0.6}^{+0.8}$ for the baryonic case $\Delta_{10} = 0.56_{-0.24}^{+0.25}$ for mesonic case. As pointed out in [$14$], the observed discrepancy in the shift of $C_9$ might arise from our incomplete understanding of the hadronic matrix elements of the two-point correlators of $\mathcal{O}_{1, \ldots, 6,8}$ with the quark electromagnetic current, which effectively shift the Wilson coefficients $C_9$ and $C_9$. This could also be due to the large experimental uncertainties for the $\Lambda_b \to \Lambda(\to p\pi^-)\mu^+\mu^-$ observables. However, if this
persists with improved statistics, this would constitute a breakdown of the universal structure of the transversity amplitudes at low recoil, as predicted by the operator product expansion.

The important distinction between the $\Lambda_b$ baryon and $B$ meson decays is the spin of the $\Lambda_b$ baryon. Therefore, the number of degrees of freedom involved in the bound state of baryon is more, hence the systematic study of $\Lambda_b \to \Lambda\gamma$ and $\Lambda_b \to \Lambda \mu^+\mu^-$ are relatively less explored in comparison to their mesonic counterpart. Also the experimental data on various $\Lambda_b$ decay channels are rather limited. Recently LHCb has reported the branching ratio of $\Lambda_b \to \Lambda\mu^+\mu^-$ [15], which is found to be lower than its SM prediction. This decay process has been extensively studied in the literature both in the SM and in various beyond the SM scenarios [16–23]. To supplement these studies, in this paper we would like to analyze the rare baryonic decay processes $\Lambda_b \to \Lambda l^+l^-$, where $l = e, \mu, \tau$ in the scalar leptoquark model. In recent times, the scalar leptoquark model has been received a lot of attention, as it can successfully explain most of the observed anomalies associated with the $b \to sll$ transitions. Leptoquarks are colour-triplet bosonic particles which can couple to a quark and lepton pair at the same time. The existence of leptoquark has been proposed in many extensions of the SM, such as grand unification model [24, 25], Pati–Salam model [26], extended technicolour model [27] and the composite models [28]. The leptoquark states can be classified as vectors (spin-1) or scalars (spin-0). They can be characterised by their Fermion no. $F = 3B + L$, where $B$ and $L$ are the baryon no. and lepton no. respectively. Scalar leptoquarks may exist at TeV scale, and can give observable signatures in various low energy processes [33]. The phenomenology of scalar leptoquarks has been studied extensively in the literature [29–37]. In this paper, we would like to study the rare baryonic decay processes $\Lambda_b \to \Lambda l^+l^-$ in the scalar leptoquark model. In particular, we estimate the decay rates, forward–backward ($A_{FB}$) and lepton polarisation asymmetries in these modes. Furthermore, we explore the possibility of lepton non-universality parameter in $\Lambda_b$ decays and also the lepton flavour violating (LFV) decays mediated via the scalar leptoquarks.

The paper is organised as follows. In section 2 we present the effective Hamiltonian responsible for the $b \to sll$ processes and the decay parameters for the semileptonic $\Lambda_b \to \Lambda l^+l^-$ decays in the SM. The NP contribution due to the exchange of scalar leptoquark has been presented in section 3 and the constraints on the leptoquark parameter space has been obtained by using the measured branching ratios of the rare decays $B_s \to l^+l^-$. In section 4, we present the numerical analysis for the branching ratios and other physical observables such as the forward–backward asymmetry, lepton polarisation asymmetry and the lepton non-universality by using the constrained leptoquark couplings. We compute the branching ratios of the LFV $\Lambda_b \to \Lambda l^+l^-$ decays in sections 5 and 6 contains the summary and conclusion.

2. Theoretical framework for the analysis of $\Lambda_b \to \Lambda l^+l^-$ decay process

In this section, we will discuss the SM contributions to the branching ratios and other physical observables of the $\Lambda_b \to \Lambda l^+l^-$, $l = e, \mu, \tau$ processes. The effective Hamiltonian describing the decay process $\Lambda_b \to \Lambda l^+l^-$ involves the quark level transition $b \to sll$ and is given by [38]

$$H_{\text{eff}} = -\frac{4G_F}{\sqrt{2}} V_{tb} V_{ts} \left[ \sum_{i=1}^{6} C_i(\mu) O_i + \frac{e}{16\pi^2} (\bar{q}_f \gamma_\mu q_i)(m_i P_L + m_b P_R) b \right] \epsilon^{\mu\nu\rho\sigma}$$

$$+ C_{\gamma e} \frac{\alpha}{4\pi} (\bar{q}_\gamma \gamma_\mu q_L) l \gamma_\nu l + C_{\gamma b} \frac{\alpha}{4\pi} (\bar{q}_\gamma \gamma_\mu q_L) l \gamma_\nu \gamma_5 l,$$

where $V_{tb}$ are the CKM matrix elements, $G_F$ denotes the Fermi constant, $\alpha$ is the fine-structure constant, $C_i$'s are the Wilson coefficients evaluated at the renormalized scale $\mu = m_b$ [39] and $P_L, P_R = (1 \mp \gamma_5)/2$ are the chiral operators. The sum over $i$ includes the current–current operators $i = 1, 2$ and the QCD–penguin operators $i = 3, 4, 5, 6$.

In addition to the short distance contributions these processes also receive additional contributions arising from the long distance effects due to the real $c\bar{c}$ resonant states of $J/\psi$, $\psi'$, i.e., $\Lambda_b \to \Lambda J/\psi(\psi') \to \Lambda l^+l^-$. These resonance contributions can be included by modifying the Wilson coefficient $C_9^\text{eff}$. Thus, the modified coefficient ($C_9^\text{eff}$) contains a perturbative part and a resonance part which can be written as

$$C_{\gamma e}^\text{eff} = C_{\gamma e}^{\text{SM}} + Y(s) + C_{\gamma e}^{\text{res}},$$

where $C_{\gamma e}^{\text{SM}}$ is the SM Wilson coefficient evaluated at the $b$ quark mass scale [39], the perturbative part $Y(s)$ receives contributions coming from one-loop matrix elements of the four quark operators [40] and the long distance resonance effect is given by

$$C_{\gamma e}^{\text{res}} = \frac{3\pi}{\alpha^2} (3C_1 + C_2 + 3C_3 + C_4 + 3C_5 + C_6) \sum_{\psi' = \psi^{(1S)}, \psi^{(2S)}, \psi^{(1D)}, \psi^{(2D)}, \psi^{(3D)}} \nu_{\psi'} \frac{m_{\psi'}}{m_b^2 - s - im_m G_\psi^{\text{eff}}},$$

2
Here the phenomenological parameter $\kappa$ is taken to be 1.65 and 2.36 [42] for the lowest resonances $J/\psi$ and $\psi'$ respectively in order to reproduce the correct branching ratio of $B(B \rightarrow J/\psi K^* \rightarrow K^* l^+ l^-) = B(B \rightarrow J/\psi K^*)B(J/\psi \rightarrow l^+ l^-)$.

The matrix elements of the hadronic currents in (1) between initial $\Lambda_b$ and the final $\Lambda$ baryon can be parameterised in terms of various form factors which are presented in appendix A. Thus, using these matrix elements, the transition amplitude for the $\Lambda_b \rightarrow \Lambda l^+ l^-$ processes can be written as [16, 18]

$$\mathcal{M}(\Lambda_b \rightarrow \Lambda l^+ l^-) = \frac{G_F}{\sqrt{2}} \frac{\alpha}{\pi} V_{tb} V_{tb}^* \left[ \hat{A}_1 \gamma^\mu \left( A_1 P_R + B_1 P_L \right) + i\sigma^{\mu\nu} q_{\nu} \left( A_2 P_R + B_2 P_L \right) \right] u_{\Lambda_b} + i \frac{G_F}{\sqrt{2}} \frac{\alpha}{\pi} V_{tb} V_{tb}^* \left[ \hat{A}_2 \gamma^\mu \left( A_1 P_R + B_1 P_L \right) + i\sigma^{\mu\nu} q_{\nu} \left( A_2 P_R + B_2 P_L \right) \right] u_{\Lambda_b},$$

where the parameters $A_i$, $B_i$, $D_j$ and $E_j$ with $i = 1, 2$, $j = 1, 2, 3$ are defined as

$$A_i = C_{i0} \frac{f_i - g_i}{2} - \frac{2m_b}{q^2} C_{i1} \frac{f_i^T + g_i^T}{2},$$
$$B_i = C_{i0} \frac{f_i + g_i}{2} - \frac{2m_b}{q^2} C_{i1} \frac{f_i^T - g_i^T}{2},$$
$$D_j = C_{j0} \frac{f_j - g_j}{2},$$
$$E_j = C_{j0} \frac{f_j + g_j}{2}.$$  \tag{5}

Using the transition amplitude (4), the double differential decay rate is given by

$$\frac{d^2 \Gamma}{d\hat{s} dz} = \frac{G_F^2 \alpha^2}{2\pi^3} \left| V_{tb} V_{tb}^* \right|^2 m_{\Lambda_b} v_1 \lambda^{1/2}(1, r, \hat{s}) K(\hat{s}, z),$$

where

$$K(\hat{s}, z) = K_{\Theta}(\hat{s}) + z K_{\Delta}(\hat{s}) + z^2 K_2(\hat{s}),$$

$$\hat{s} = s/m_{\Lambda_b}^2 \text{ and } z = \hat{p}_b \cdot \hat{P}_p^+ \text{ is the angle between the momenta of } \Lambda_b \text{ and } l^+ \text{ in the dilepton invariant mass frame.}$$

The complete expressions for $K_{\Theta}(\hat{s})$, $K_{\Delta}(\hat{s})$ and $K_2(\hat{s})$ are given in appendix B. Here $v_1 = \sqrt{1 - (4m_l^2/q^2)}$ and $\lambda(1, r, \hat{s}) = (1 - r^2 - 2\hat{s}(1 + r) + \hat{s}^2) \text{ is the triangle function with } r = m_{\Lambda_b}/m_{\Lambda_b}$.

The physical allowed range for $s \equiv q^2$ is

$$4m_l^2 \leq s \leq (m_{\Lambda_b} - m_{\Lambda_b})^2.$$  \tag{8}

Another interesting observable is the zero-crossing of the forward–backward asymmetry, wherein the position of the zero value of the forward–backward asymmetry parameter ($A_{FB}$) is very useful to look for the NP signal.

The normalised forward–backward asymmetry is defined as

$$A_{FB}(\hat{s}) = \frac{\int_0^1 \frac{d\Gamma}{d\hat{s} dz} dz - \int_{-1}^0 \frac{d\Gamma}{d\hat{s} dz} dz}{\int_0^1 \frac{d\Gamma}{d\hat{s} dz} dz + \int_{-1}^0 \frac{d\Gamma}{d\hat{s} dz} dz},$$

which can be simplified to

$$A_{FB}(\hat{s}) = \frac{K_{\Delta}(\hat{s})}{K_{\Theta}(\hat{s}) + K_2(\hat{s})/3}.$$  \tag{10}

The polarisation asymmetries $P_i (i = L, N, T)$ are defined as

$$P_i(\hat{s}) = \frac{d\Gamma}{d\hat{s}}(\hat{s} = \hat{e}_i) - \frac{d\Gamma}{d\hat{s}}(\hat{s} = -\hat{e}_i),$$

where $\hat{e}_i$'s are the unit vectors along the longitudinal, normal and transverse components of the $l^+$ polarisation and $\hat{s}$ is a unit vector, used to write the $l^+$ four-spin vector $(s_\perp)$, along the $l^+$ spin in its rest frame as

$$s_\perp = \hat{p} \times \frac{\hat{e}}{m_l}, \quad s_\perp = \hat{s} + \frac{\hat{e}}{E_{l^+} + m_l} \hat{p}.$$  \tag{12}

Thus, the observables $P_L$, $P_T$ and $P_N$ correspond to longitudinal, transverse and normal polarisation asymmetries respectively. The observables $P_L$ and $P_T$ are $P$-odd, $T$-even, while $P_N$ is $P$-even, $T$-odd and CP-odd. The explicit expressions for forward–backward asymmetry and all the polarisation parameters are taken from [16, 18, 19].
Another interesting observable is the lepton universality violation (LUV) parameter, which has been recently observed by the LHCb collaboration in \( B^+ \rightarrow K^{*+}l^+l^- \) process and has 2.6\( \sigma \) deviation from its SM predicted value [6]. Analogously, we define the parameter \( (R_\lambda) \) as the ratio of branching fractions of \( \Lambda_b \rightarrow \Lambda l^+l^- \) into dimuon over dielectron as

\[
R_\lambda = \frac{\text{Br}(\Lambda_b \rightarrow \Lambda \mu^+\mu^-)}{\text{Br}(\Lambda_b \rightarrow \Lambda e^+e^-)}.
\]

(13)

3. NP contribution due to scalar leptoquark exchange

In this section we will consider the effect of scalar leptoquarks to the \( \Lambda_b \rightarrow \Lambda l^+l^- \) decay processes. The exchange of leptoquarks will contribute additional operators to the SM effective Hamiltonian and thus, the various observables may deviate significantly from their corresponding SM values. The scalar leptoquark multiplets with representations \( X (3, 2, 7/6) \) and \( X (3, 2, 1/6) \) under the SM gauge group \( SU(3)_C \times SU(2)_L \times U(1)_Y \) conserve baryon and lepton numbers and do not allow proton decay. These baryon and lepton number conserving scalar leptoquarks can have sizable Yukawa couplings and could be light enough to be accessible in accelerator searches. Thus, they could potentially contribute to the \( b \rightarrow s l^+l^- \) transitions and one can constrain the underlying couplings from experimental data on \( B_s \rightarrow l^+l^- \) processes as well as from \( B_s \rightarrow \bar{B}_s \) mixing.

The interaction Lagrangian of the scalar leptoquarks \( X = (3, 2, 7/6) \) with the SM bilinear fermions is given as [33, 34]

\[
\mathcal{L} = -\lambda_{ij}^L \bar{u}_i X L_j^L - \lambda_{ij}^R \bar{e}_i X Q_j^L + \text{h.c.},
\]

(14)

where \( i, j \) are the generation indices, \( X \) is the leptoquark doublet, \( Q_L (L_L) \) denotes the left handed quark (lepton) doublet, the right handed up-type quark (charged lepton) singlet is represented by \( u_R (\bar{e}_R) \) and \( \epsilon = i\sigma_2 \) is a \( 2 \times 2 \) matrix. The multiplets defined above are represented as

\[
X = \begin{pmatrix} V_\alpha \\ Y_\alpha \end{pmatrix}, \quad Q_L = \begin{pmatrix} u_L \\ d_L \end{pmatrix}, \quad \text{and} \quad L_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix}.
\]

(15)

Now expanding the \( SU(2) \) indices, the interaction Lagrangian (14) takes the form

\[
\mathcal{L} = -\lambda_{ij}^{L\alpha} \bar{u}^{\alpha}_{i \lambda} (V_\alpha e^\dagger_j - Y_{\alpha} \nu^j_L) - \lambda_{ij}^{R\alpha} \bar{e}^{\alpha}_{i \lambda} (V^\dagger_\alpha u_{j \lambda} + Y_{\alpha} d_{j \lambda}^\dagger) + \text{h.c.}.
\]

(16)

Thus, from equation (16) one can obtain the interaction Hamiltonian for \( b \rightarrow s l^+l^- \) processes after performing the Fierz transformation as

\[
\mathcal{H}_{LQ} = \frac{\lambda_{ij}^{L\alpha} \lambda_{ij}^{R\alpha*}}{8M_W^2} [\bar{b} \gamma^\mu (1 - \gamma_5) b][\bar{l}_i \gamma_\mu (1 + \gamma_5) l_j] = \frac{\lambda_{ij}^{L\alpha} \lambda_{ij}^{R\alpha*}}{4M_W^2} (O_9 + O_{10}).
\]

(17)

Comparing (17) with the corresponding SM effective Hamiltonian (1), one can obtain the new Wilson coefficients as

\[
C_{i\alpha}^{NP} = C_{i\alpha}^{NP} = -\frac{\pi}{2\sqrt{2} G_F \alpha} \frac{\lambda_{ij}^{L\alpha} \lambda_{ij}^{R\alpha*}}{M_W^2}.
\]

(18)

Similarly, the interaction Lagrangian due to the exchange of the scalar leptoquark \( X = (3, 2, 1/6) \) is

\[
\mathcal{L} = -\lambda_{ij}^R \bar{d}_{i \lambda} (V_\alpha e^\dagger_j - Y_{\alpha} \nu^j_L) + \text{h.c.},
\]

(19)

which contributes to the primed Wilson coefficients \( C_{i\alpha}^{NP} \) corresponding to the semileptonic electroweak penguin operators \( O_{9,10}' \) (i.e., the right-handed counter parts of the SM operators \( O_{9,10} \)) and are given as

\[
C_{i\alpha}^{NP} = -C_{i\alpha}^{NP} = -\frac{\pi}{2\sqrt{2} G_F \alpha} \frac{\lambda_{ij}^R \lambda_{ij}^{R\alpha*}}{M_W^2}.
\]

(20)

Thus, from the above equations (18) and (20), one can find that there are four additional Wilson coefficients \( C_{i\alpha}^{NP} \), which will contribute to the \( b \rightarrow s l^+l^- \) processes due to the scalar leptoquark exchange. Thus, the modified parameters (5) in the amplitude (4), become
The range of the experimental data is \(|\Delta f|\), where \(f\) is the SM branching ratio and the parameters \(f_{NP}\) are defined as

\[
\begin{align*}
A_i &= C_{ij}^{NP} \frac{g_i^2}{2} + \left( C_{ij}^{eff} + C_{ij}^{NP} \right) \frac{f_i - g_i}{2} - \frac{2m_{t_i} f_{SM}}{q^2} C_{ij}^{SM} \frac{f_i^T + g_i^T}{2}, \\
B_i &= \left( C_{ij}^{eff} + C_{ij}^{NP} \right) \frac{g_i}{2} - \frac{2m_{t_i} f_{SM}}{q^2} C_{ij}^{SM} \frac{f_i^T - g_i^T}{2} + C_{ij}^{NP} \frac{f_i - g_i}{2}, \\
D_i &= C_{ij}^{NP} \frac{f_i + g_i}{2} + \left( C_{ij}^{16} + C_{ij}^{NP} \right) \frac{g_i}{2}, \\
E_i &= \left( C_{ij}^{SM} + C_{ij}^{16} \right) \frac{g_i}{2} + C_{ij}^{NP} \frac{f_i - g_i}{2}.
\end{align*}
\] 

Next, we have to find out the constraints on the leptoquark couplings to see how various observables behave in the LQ model. The detailed calculation of the constraint on the new leptoquark parameter space has been presented in [29–31], therefore here we will simply quote the main results. We constrain the leptoquark coupling using the equations

\[
\text{Bound on the LQ couplings (GeV$^{-2}$)}
\]

from the SM theoretical prediction of \(B_{i \to j} \to l^+l^-\) process and \(B_{i \to j} \to l^+l^-\) mixing, however, the technical details can be found in [29–31].

### 3.1. Constraint from \(B_{i \to j} \to \mu^+\mu^-\) process

In the leptoquark model the branching ratio for the \(B_{i \to j} \to \mu^+\mu^-\) mode can be given as

\[
\text{Br}(B_{i \to j} \to \mu^+\mu^-) = \frac{G_F^2}{16\pi^3} \alpha^2 f_{\mu}^2 m_{\mu}^2 \left| V_{ik} V_{jk}^* \right| C_{ij}^{NP} f_{SM}^{NP} + C_{ij}^{NP} f_{SM}^{NP} f_{SM}^{NP} \left[ 1 - \frac{4m_{\mu}^2}{M_{\mu}^2} \right],
\]

where \(\text{Br}^{SM}\) is the SM branching ratio and the parameters \(r\) and \(\phi^{NP}\) are defined as

\[
\begin{align*}
\text{Br}^{NP} &= \frac{C_{ij}^{NP} - C_{ij}^{SM}}{C_{ij}^{SM}}, \\
\text{Br}^{NP} &= \frac{C_{ij}^{NP} - C_{ij}^{SM}}{C_{ij}^{SM}}.
\end{align*}
\]

Now comparing the SM theoretical prediction of \(\text{Br}(B_{i \to j} \to \mu^+\mu^-)\) [44]

\[
\text{Br}(B_{i \to j} \to \mu^+\mu^-)_{SM} = (3.65 \pm 0.23) \times 10^{-9},
\]

with the corresponding experimental value [45–47]

\[
\text{Br}(B_{i \to j} \to \mu^+\mu^-) = (2.9 \pm 0.7) \times 10^{-9},
\]

and assuming that each individual leptoquark contribution to the branching ratio does not exceed the experimental result, one can obtain the bound on the NP parameters \(r\) and \(\phi^{NP}\). The allowed parameter space in \(r - \phi^{NP}\) plane which is compatible with the 1 \(\sigma\) range of the experimental data is

\[
0 \leq r \leq 0.35, \quad \text{with} \quad \pi/2 \leq \phi^{NP} \leq 3\pi/2.
\]

These bounds can be translated to obtain the bounds for the leptoquark couplings as

\[
0 \leq \frac{\left| \lambda_{ij}^{\mu} \lambda_{ij}^{\nu} \right|}{M_{L}^2} = \frac{\left| \lambda_{ij}^{\mu} \lambda_{ij}^{\nu} \right|}{M_{L}^2} \leq 5 \times 10^{-9} \text{ GeV}^{-2} \quad \text{for} \quad \pi/2 \leq \phi^{NP} \leq 3\pi/2.
\]

Similarly, one can obtain the upper bound on the product of various combination of leptoquark couplings from \(B_{i \to j} \to l^+l^-\) processes which are presented in table 1. Using the bounds on leptoquark couplings one can obtain the constraints on new Wilson coefficients using the equations (18) and (20).

### Table 1. Constraints on the leptoquark couplings obtained from various leptonic \(B_{i \to j} \to l^+l^-\) decays [29], where \(M_L\) denotes the mass of the scalar LQ.

| Decay Process | Couplings involved | Bound on the LQ couplings (GeV$^{-2}$) |
|---------------|--------------------|--------------------------------------|
| \(B_i \to \mu^+\mu^-\) | \(\frac{1}{M_{L}^2} g_i^2 \) | \(< 5.0 \times 10^{-9}\) |
| \(B_i \to e^+e^-\) | \(\frac{1}{M_{L}^2} g_i^2 \) | \(< 2.54 \times 10^{-5}\) |
| \(B_i \to \tau^+\tau^-\) | \(\frac{1}{M_{L}^2} g_i^2 \) | \(< 1.2 \times 10^{-8}\) |
3.2. Constraint from $B_s - \bar{B}_s$ mixing

In this subsection, we will discuss the constraint on leptoquark couplings from the $B_s - \bar{B}_s$ mixing, which in the SM, proceeds through the box diagram with internal top quark and $W$ boson exchange. The effective Hamiltonian describing the $\Delta B = 2$ transition is given as [48]

$$\mathcal{H}_{\text{eff}} = \frac{G_F^2}{16\pi^2} V_{tb} V^*_{t\bar{s}} M^2_{B_s} S_0(x_i) \eta_f \langle \bar{s} b \rangle V_{\nu - A} \langle \bar{s} \bar{b} \rangle V_{\nu - A},$$

(28)

where $\eta_f$ is the QCD correction factor and $S_0(x_i)$ is the loop function given in [48].

Thus, the $B_s - \bar{B}_s$ mixing amplitude in the SM, can be written as

$$M^2_{2s} = \frac{1}{2M_{B_s}} \langle \bar{B}_s \rangle \mathcal{H}_{\text{eff}} \langle B_s \rangle = \frac{G_F^2}{12\pi^2} M^2_{B_s} |V_{tb} V^*_{t\bar{s}}|^2 \eta_f \bar{B}_f f^2 \mu M_{B_s} S_0(x_i).$$

(29)

The corresponding mass difference can be computed from the mixing amplitude through $\Delta M_s = 2|M_{2s}|$. Now using the particle masses from [8], $\eta_f = 0.551$, the Bag parameter $\sqrt{B_{B_s}} = 1.320 \pm 0.017 \pm 0.030$ and the decay constant $f_{B_s} = 225.6 \pm 1.1 \pm 5.4$ from [49], the value of $\Delta M_s$ in the SM, is found as

$$\Delta M^2_{s} = (17.426 \pm 1.057) \text{ ps}^{-1},$$

(30)

which is in good agreement with the experimental result [8]

$$\Delta M_s = 17.761 \pm 0.022 \text{ ps}^{-1}. \quad (31)$$

For $X(3, 2, 7/6)$ LQ, the mixing amplitude receives additional contribution from leptoquark and charged lepton in the box diagram whereas for $X(3, 2, 1/6)$ both charged lepton and neutrino will contribute to the mixing amplitude. The effective Hamiltonian due to the leptoquark $X(3, 2, 7/6)$ is given by

$$\mathcal{H}_{\text{eff}} = \sum_{i=e,\mu,\tau} \frac{(\lambda^{i\bar{q}}\lambda^{\bar{q}q})^2}{128\pi^2} \left[ \frac{1}{M^2_S} \left( \bar{B}_f f^2 \mu M_{B_s} \right) \left( \bar{B}_{\gamma^iP_L s} \right) \left( \bar{B}_{\gamma^iP_L s} \right) \right],$$

(32)

and for $X(3, 2, 1/6)$ leptoquark the corresponding effective Hamiltonian becomes

$$\mathcal{H}_{\text{eff}} = \sum_{i=e,\mu,\tau} \frac{(\lambda^{i\bar{q}}\lambda^{\bar{q}q})^2}{128\pi^2} \left[ \frac{1}{M^2_S} \left( \bar{B}_f f^2 \mu M_{B_s} \right) \left( \bar{B}_{\gamma^iP_R s} \right) \left( \bar{B}_{\gamma^iP_R s} \right) \right],$$

(33)

where the loop function $I(x)$ is given as

$$I(x) = \frac{1 - x^2 + 2x \log x}{(1 - x)^2}. \quad (34)$$

Thus, the contribution to the mixing amplitude due to the exchange of scalar leptoquark is given by

$$M^2_{2s}^{\text{LQ}} = \frac{(\lambda^{i\bar{q}}\lambda^{\bar{q}q})^2}{128\pi^2} \left[ \frac{1}{M^2_S} \left( \bar{B}_f f^2 \mu M_{B_s} \right) \right], \quad \text{for } X(3, 2, 1/6), \quad (35)$$

$$M^2_{2s}^{\text{LQ}} = \frac{(\lambda^{i\bar{q}}\lambda^{\bar{q}q})^2}{128\pi^2} \left[ \frac{1}{M^2_S} \left( \bar{B}_f f^2 \mu M_{B_s} \right) \right], \quad \text{for } X(3, 2, 7/6). \quad (36)$$

Including both the SM and leptoquark contributions the total mass difference is given as

$$\Delta M_s = \Delta M^2_{s}^{\text{SM}} \left[ 1 + \frac{c}{16G_F^2 |V_{tb} V^*_{t\bar{s}}|^2 \eta_f S_0(x_i)} \left( \frac{\lambda^{i\bar{q}}\lambda^{\bar{q}q}}{M^2_S} \right) \right],$$

(36)

where the constant $c = 1$ for $X(3, 2, 1/6)$ and $1/2$ for $X(3, 2, 7/6)$. Now varying the mass difference ($\Delta M_s / \Delta M^2_{s}^{\text{SM}}$) within its 1σ allowed range [8], the constraint on $|$X^2\lambda^{i\bar{q}}\lambda^{\bar{q}q}|/M_S|$ is found to be [31]

$$0 \leq \frac{\lambda^{i\bar{q}}\lambda^{\bar{q}q}}{M_S} \leq 7.5 \times 10^{-5} \text{ GeV}^{-1}, \quad \text{for } X(3, 2, 7/6),$$

$$0 \leq \frac{\lambda^{i\bar{q}}\lambda^{\bar{q}q}}{M_S} \leq 5.0 \times 10^{-5} \text{ GeV}^{-1}, \quad \text{for } X(3, 2, 1/6). \quad (37)$$

In order to relate this results with the bounds obtained $B_s \rightarrow \mu \mu$ process, we scale the couplings obtained from $B_s - \bar{B}_s$ mass difference for a benchmark leptoquark mass of 1 TeV and the bounds in equation (37) is translated as
which are reasonably higher than those of obtained from $B_s \to \mu\mu$ process. Hence in our analysis, we will use the bounds (26) as discussed in the previous subsection.

4. Numerical analysis

After having the detailed knowledge about the SM observables and the bound on the new leptoquark couplings, we now proceed for numerical analysis. We have taken the particle masses and the life time of $\Lambda_b$ baryon from [8]. The $q^2$ dependence of form factors derived in the light cone sum rule (LCSR) approach can be parameterised as

$$f_i(q^2) = \frac{f_i(0)}{1 - a(q^2/m_{\Lambda_b}^2) + b(q^2/m_{\Lambda_b}^2)^2},$$

where the values of the parameters $f_i(0), a$ and $b$ and are listed in table 2 [21]. The other form factors are related to these two and the HQET form factors ($F_{i\gamma}$) through [21]

$$F_i = g_i = 1 = F_i + \frac{m_{\Lambda_b}}{m_{\Lambda_b}},$$

$$f_2 = g_2 = F_2 - \frac{m_{\Lambda_b}}{m_{\Lambda_b}}.$$  

In the lattice QCD formalism, the $\Lambda_b \to \Lambda$ helicity form factors, i.e., $f_{\pm,\pm,0}^T, g_{\pm,\pm,0}^T, h_{\pm,\pm,0}^T$ and $h_{\pm,\pm,0}^T$ in the physical limit can have the simple form [23]

$$f(q^2) = \frac{1}{1 - q^2/(m_{\Lambda_b}^2)} \left[ a_0^f + a_1^f [z(q^2)] + a_2^f [z(q^2)]^2 \right],$$

where the values and uncertainties of the parameters $\alpha_0, \alpha_1$ and $\alpha_2$ from the higher-order fit are given in table 5 of [23]. These helicity form factors are related to the form factors $F_i^{(T)}$ and $g_i^{(T)}$ used in this work as follows:

$$f_+ = f_1 - \frac{q^2}{m_{\Lambda_b}^2} f_2, \quad f_- = f_1 - (m_{\Lambda_b} + m_{\Lambda_b}) f_2, \quad f_0 = f_1 + \frac{q^2}{m_{\Lambda_b} - m_{\Lambda_b}} f_2,$$

$$g_+ = g_1 + \frac{q^2}{m_{\Lambda_b}^2} g_2, \quad g_- = g_1 + (m_{\Lambda_b} + m_{\Lambda_b}) g_2, \quad g_0 = g_1 - \frac{q^2}{m_{\Lambda_b} + m_{\Lambda_b}} g_2,$$

$$h_+ = f_2^T - \frac{m_{\Lambda_b} + m_{\Lambda_b}}{q^2} f_2^T, \quad h_- = f_2^T - \frac{m_{\Lambda_b} + m_{\Lambda_b}}{m_{\Lambda_b} + m_{\Lambda_b}} f_2^T,$$

$$h_0 = g_2^T + \frac{m_{\Lambda_b} - m_{\Lambda_b}}{q^2} g_2^T, \quad h_0 = g_2^T + \frac{m_{\Lambda_b} - m_{\Lambda_b}}{m_{\Lambda_b} - m_{\Lambda_b}} g_2^T.$$  

In our analysis, we have taken the form factors computed in the LCSR approach for low $q^2$ region (as these are not so well-behaved in the high $q^2$ regime), and for high $q^2$ theory we have used the lattice QCD calculations of $\Lambda_b \to \Lambda$ form factors [23]. The values of the Wilson coefficients used in our analysis are evaluated at the
renormalisation scale $\mu \approx m_b = 4.8$ GeV. In the LQ model, the NP contributions to the branching ratios and forward-backward asymmetry parameters are encoded in the new Wilson coefficients. By using the above input parameters and the values of the new Wilson coefficients, we show in figure 1, the $q^2$ variation of branching ratio of $\Lambda_b \rightarrow \Lambda e^+e^-$ (top left panel), $\Lambda_b \rightarrow \Lambda \mu^+\mu^-$ (top right panel) and $\Lambda_b \rightarrow \Lambda \tau^+\tau^-$ (bottom panel) for the full kinematically accessible physical region. In these plots, we have shown the contributions arising from the exchange of $X = (3, 2, 7/6)$ leptoquark model. The SM contributions are represented by blue lines and the grey bands denote the theoretical uncertainties arising due to the uncertainties associated with the CKM matrix elements and the hadronic form factors. The green bands represent the leptoquark contributions to the branching ratios. The bin-wise experimental results for $\Lambda_b \rightarrow \Lambda \mu^+\mu^-$ process [15] are shown by black data points. There is slight deviation in the decay distribution between the predicted and observed data. The corresponding results coming from the exchange of the $X = (3, 2, 1/6)$ LQ are shown in figure 2. From these figures, one can see that the branching ratios of $\Lambda_b \rightarrow \Lambda e^+e^-$ and $\Lambda_b \rightarrow \Lambda \tau^+\tau^-$ decay processes deviate significantly from their SM predictions, whereas the NP effects on $\Lambda_b \rightarrow \Lambda \mu^+\mu^-$ branching ratio is not so prominent. In table 3, we present the integrated values of branching ratio for all the above processes, where we have used the veto windows as $(8 \text{ GeV}^2 < m_{f^+}^2 < 11 \text{ GeV}^2)$ and $(12.5 \text{ GeV}^2 < m_{f^+}^2 < 15 \text{ GeV}^2)$ [15], to eliminate the backgrounds coming from the dominant resonances $\Lambda_b \rightarrow J/\psi(\psi')$ with $J/\psi(\psi') \rightarrow l^+l^-$. The predicted branching ratio for $\Lambda_b \rightarrow \Lambda \mu^+\mu^-$ is almost consistent with the observed data $\text{Br}(\Lambda_b \rightarrow \Lambda \mu^+\mu^-) = (1.08 \pm 0.28) \times 10^{-6}$ [8]. Also, as seen from table 3, the experimental result can be accommodated in the leptoquark model. Within the SM, the forward backward asymmetry parameters in the $B \rightarrow Kl^+l^-$ decay processes are identically zero since they only involve scalar and tensor types of currents, whereas $B \rightarrow Kl^+l^-$ processes are described by only vector-type interactions. However, for semileptonic $\Delta_m \rightarrow \Lambda l^+l^-$ decay processes, the forward backward asymmetry depends on two combinations of the Wilson coefficients $\text{Re}(C_{17}^\text{eff} C_{18}^\mu)\text{ and Re}(C_{19}^\text{eff} C_{19}^\mu)$ [16] and thus, can have negative values in the SM. The contribution due to the new Wilson coefficients ($C_{9,10}^\text{NP}$) may enhance the rate of asymmetries and can shift the zero position of these asymmetries. In figure 3, the variation of forward-backward asymmetry for $\Lambda_b \rightarrow \Lambda \mu^+\mu^-$ (left panel), $\Lambda_b \rightarrow \Lambda \tau^+\tau^-$ (right panel) modes are depicted with respect to $q^2$ both in the SM and in the $X = (3, 2, 7/6)$ LQ model including the LD contributions and the corresponding integrated values are presented in table 3. Similarly the variation of forward-backward asymmetry for $X = (3, 2, 1/6)$ LQ exchange are shown in figure 4. We found no significant deviation of the zero position of $\Lambda_B$ from its SM value due to the leptoquark contributions in $\Lambda_b \rightarrow \Lambda \mu^+\mu^-$ process. However, there is certain discrepancy between the observed and predicted results in
the high $q^2$ regime. The forward–backward asymmetry for $\Lambda_b \to \Lambda \tau^+ \tau^-$ process however, has significant deviation from the SM in both the $X = (3, 2, 7/6)$ and $X = (3, 2, 1/6)$ leptoquark model. Besides the branching ratios and forward–backward asymmetry parameters of $\Lambda_b \to \Lambda \mu^+ \mu^-$ processes, the NP effects can also be observed in the lepton polarisation asymmetries. In the left panel of figure 5, the distribution of the longitudinal (top), transverse (middle) and normal (bottom) polarisation components for $\Lambda_b \to \Lambda \tau^+ \tau^-$ process are shown both in the SM and in the $X = (3, 2, 7/6)$ leptoquark model, and the corresponding plots for $\Lambda_b \to \Lambda \mu^+ \mu^-$ process are presented in the right panel. The integrated values of all the three polarisations in the full physical phase space have been presented in table 3. In figure 6, we have shown the variation of the different polarisation parameters for $\Lambda_b \to \Lambda \mu^+ \mu^-$ process in the $X = (3, 2, 1/6)$ leptoquark model. It is found from table 3, that the transverse and normal polarisation values are very small in the SM and even the leptoquark model does not give any significant deviation.

**Figure 2.** Same as figure 1 for $X = (3, 2, 1/6)$ LQ exchange.

**Table 3.** The predicted integrated values of the branching ratio, forward–backward asymmetry, lepton polarisation asymmetry and the lepton non-universality with respect to their respective $q^2$ range for the $\Lambda_b \to \Lambda \mu^+ \tau^- \mu^-$ processes in the SM and the LQ model.

| Observables | SM prediction | Values in $Y = 7/6$ LQ model | Values in $Y = 1/6$ LQ model |
|-------------|---------------|-----------------------------|-----------------------------|
| $\text{Br}(\Lambda_b \to \Lambda e^+ e^-)$ | $(1.168 \pm 0.134) \times 10^{-6}$ | $(1.168 - 1.91) \times 10^{-6}$ | $(1.168 - 2.13) \times 10^{-6}$ |
| $\text{Br}(\Lambda_b \to \Lambda \mu^+ \mu^-)$ | $(1.165 \pm 0.132) \times 10^{-6}$ | $(1.165 - 1.37) \times 10^{-6}$ | $(1.165 - 1.52) \times 10^{-6}$ |
| $\langle A_{FB}^L \rangle$ | $-0.567$ | $-0.567 \to -0.446$ | $-0.567 \to -0.54$ |
| $\langle P_T^L \rangle$ | $0.34$ | $0.3 \to 0.34$ | $0.24 \to 0.34$ |
| $\langle P_T^T \rangle$ | $-4.5 \times 10^{-4}$ | $-(4.5 \to 2.87) \times 10^{-4}$ | $-(0.45 \to 3.26) \times 10^{-3}$ |
| $\langle P_T^N \rangle$ | $-0.0192$ | $-0.0192 \to -0.013$ | $-0.0192 \to -0.012$ |
| $\text{Br}(\Lambda_b \to \Lambda \tau^+ \tau^-)$ | $(2.13 \pm 0.215) \times 10^{-7}$ | $(2.13 - 4.38) \times 10^{-7}$ | $(2.13 - 8.32) \times 10^{-7}$ |
| $\langle A_{FB}^L \rangle$ | $-0.38$ | $-0.38 \to 3.2 \times 10^{-3}$ | $-0.38 \to 7.68 \times 10^{-2}$ |
| $\langle P_T^L \rangle$ | $0.075$ | $0.047 \to 0.075$ | $6.3 \times 10^{-3} \to 0.075$ |
| $\langle P_T^T \rangle$ | $-2.3 \times 10^{-3}$ | $-(7 \to 2.3) \times 10^{-3}$ | $-(0.23 \to 2.0) \times 10^{-2}$ |
| $\langle P_T^N \rangle$ | $-0.05$ | $-0.05 \to 8.1 \times 10^{-3}$ | $-0.05 \to 0.0316$ |
| $\langle R_{\tau^+ \tau^-}^e \rangle$ | 0.997 | $0.67 \to 0.997$ | $0.68 \to 0.997$ |
| $\langle R_{\tau^+ \tau^-}^\mu \rangle_{L \leq (1.62)}$ | 0.998 | $0.67 \to 0.998$ | $0.74 \to 0.998$ |
Analogous to the lepton flavour non-universality parameter $R_K$, i.e., the ratio of branching fractions of $B\rightarrow K\mu^+\mu^-$ over $B\rightarrow Ke^+e^-$, we would like to see whether it is possible to observe lepton non-universality in the $\Lambda_b$ decays. We have defined these parameters as e.g., $R_L^{\mu e} = \text{Br}(\Lambda_b \rightarrow \Lambda \mu^+\mu^-)/\text{Br}(\Lambda_b \rightarrow \Lambda e^+e^-)$. In figure 7, we show the variation of lepton nonuniversality parameter $R_L^{\mu e}$ (top-right panel), $R_L^{e\tau}$ (bottom-left panel) and $R_L^{\mu\tau}$ (bottom-right panel) in their respective $q^2$ region. Also, we show the low-$q^2$ behaviour of $R_L^{\mu e}$ (top-left panel), in the range $1 \leq q^2 \leq 6 \text{ GeV}^2$. These results are for $X = (3, 2, 7/6)$ leptoquark. Similarly the lepton nonuniversality plot for $X = (3, 2, 1/6)$ leptoquark exchange is shown in figure 8. The integrated values of the lepton non-universality parameter in both SM and LQ model are presented in table 3. We found that there is significant violation of lepton universality in $\Lambda_b$ decays, though there is no experimental evidence so far. The violation of lepton universality is more pronounced for the processes having $\tau$ as final particle. However, as the reconstruction of tau events are extremely difficult, this observable may not be sensitive enough to be observed in near future. As seen from the top-left panel of figures 7 and 8, the parameter $R_L^{\mu e}$ is very promising for the Belle II experiment, as the LHCb, being a hadronic machine works better in muon mode than electron.

5. LFV $\Lambda_b \rightarrow \Lambda l_l^-l_l^+$ decays

In this section, we will compute the branching ratios of LFV $\Lambda_b$ decays mediating through the exchange of scalar leptoquarks. The LFV decay processes are extremely rare in the SM as they are either two-loop suppressed with tiny neutrino masses in one of the loop or proceed through box diagram (which is also highly suppressed due to tiny neutrino mass). However, they can occur at tree level in the LQ model and are expected to have significantly large branching fractions. The observation of neutrino oscillation has provided unambiguous evidence for lepton flavour violation in the neutral lepton sector which in turn provides motivation to explore other LFV transitions such as $l_l \rightarrow l_l^-\gamma, l_l^- \rightarrow l_l^-l_l^-l_l^-, B \rightarrow l_l^+l_l^-$ etc. Though there is no direct experimental evidence for such processes, but there exists experimental upper bounds on some of these modes. The LFV decays in the $B$ meson and in the charged lepton sector have been widely investigated in the literature \cite{29, 36, 50}. Therefore, it is interesting to see whether LFV decays could be observed in $\Lambda_b$ decays also.
As discussed earlier, these processes occur at tree level due to the exchange of scalar leptoquarks. In the leptoquark model the effective Hamiltonian for \( \Lambda_b \rightarrow \Lambda_l^- \mu^+ \ell^- \) process is given as \( H_{\text{LQ}} \) as \( \chi^3 \gamma \mu P_L (\gamma \mu (1 + \gamma_5)) I_j \) where the coefficient \( G_{\text{LQ}} \) is
\[
G_{\text{LQ}} = \frac{\chi^3 \lambda_{ij}^{2\pi}}{8M_W^2}.
\]
Using the form factors given in the appendix A, the amplitude for the LFV \( \Lambda_b \rightarrow \Lambda_l^- \mu^+ \ell^- \) decay is given by
\[
\mathcal{M}(\Lambda_b \rightarrow \Lambda_l^- \mu^+ \ell^-) = G_{\text{LQ}} (\gamma \mu (1 + \gamma_5) I_j) \left( a_{ij} (\nu \tau (\nu \tau P_R + B_P (P_L)) u_{\Lambda_b} + \bar{u}_{\Lambda_b} \tau \mu (\nu \tau P_R + B_P (P_L)) u_{\Lambda_b} + q^\tau \bar{u}_{\Lambda_b} (\nu \tau P_R + B_P (P_L)) u_{\Lambda_b} \right).
\]
The coefficients \( A_k' \) and \( B_k' \) in (45) are related to the form factors through
\[
A_k' = \frac{f_k - g_k}{2} \quad \text{and} \quad B_k' = \frac{f_k + g_k}{2}, \quad k = 1, 2, 3.
\]
Now using this transition amplitude, the branching ratio for the \( \Lambda_b \rightarrow \Lambda_l^- \mu^+ \ell^- \) process is given as
\[
\frac{d^3 \Gamma}{d\xi d\cos \theta} = \frac{|G_{\text{LQ}}|^2 m_{\Lambda_b}}{2\pi^3} \frac{\sqrt{M_{\Lambda_b} \lambda_2}}{\xi} I(\xi),
\]

Figure 5. The plots in the left panel represent the longitudinal (top), transverse (middle) and normal (bottom) polarisations for \( \Lambda_b \rightarrow \Lambda_l^- \mu^+ \ell^- \) process with respect to \( q^3/m_{\Lambda_b} \) in the \( X = (3, 2, 7/6) \) LQ model. The corresponding plots for \( \Lambda_b \rightarrow \Lambda \tau^+ \tau^- \) mode are shown in the right panel.
where

\[ I(\hat{s}) = I_0(\hat{s}) + I_1(\hat{s}) \cos \theta + I_2(\hat{s}) \cos^2 \theta, \]

(48)
\[ I_0(\hat{s}) = \frac{1}{4} (\vert A_1^2 \vert^2 + \vert B_1^2 \vert^2 + m_{\lambda s}^2 \delta (\vert A_2^2 \vert^2 + \vert B_2^2 \vert^2))(1 - r^2 - \hat{s}^2) \]
\[ - 2 \sqrt{\hat{s}} \left( 1 - \frac{m_1^2 + m_2^2}{q^2} \right) \left( \text{Re}(A_1^* B_1^{* \prime}) + m_{\lambda s} \delta \text{Re}(A_2^* B_2^{* \prime}) \right) \]
\[ - \text{Re}(A_2^* B_1^{* \prime}) \sqrt{\hat{r}} \left[ - \frac{\hat{s}(m_1^2 - m_2^2)^2}{m_{\lambda s}^2} + \hat{s}(m_1^2 + m_2^2) \right] \]
\[ + 2m_{\lambda s} \delta \left( 1 - \frac{m_1^2 + m_2^2}{2q^2} \right) \left[ (\text{Re}(A_1^* A_2^{* \prime}) + \text{Re}(B_1^* B_2^{* \prime})) \sqrt{\hat{r}} (1 - t) \right. \]
\[ - (\text{Re}(A_1^* B_2^{* \prime}) + \text{Re}(B_1^* A_2^{* \prime}))(t - r) \]
\[ + \frac{(m_1^2 + m_2^2)^2}{m_{\lambda s}^2} \left[ (\text{Re}(A_1^* A_2^{* \prime}) + \text{Re}(B_1^* B_2^{* \prime})) \sqrt{\hat{r}} (1 - t) + (\text{Re}(A_1^* B_2^{* \prime}) + \text{Re}(B_1^* A_2^{* \prime}))(t - r) \right] \]
\[ + \left[ \hat{s}(m_1^2 + m_2^2) - \frac{(m_1^2 - m_2^2)^2}{m_{\lambda s}^2} \left[ \frac{t}{2} (\vert A_1^2 \vert^2 + \vert B_1^2 \vert^2) - \sqrt{\hat{r}} \text{Re}(A_1^* B_2^{* \prime}) \right] \right], \] (49)

\[ \hat{h}(\hat{s}) = \frac{\sqrt{\lambda s}}{\hat{s}} \left[ - \frac{1}{2} \hat{s}(\vert A_1^2 \vert^2 - \vert B_1^2 \vert^2 + (m_1^2 - m_2^2)(1 - t - \hat{s})(\vert A_1^2 \vert^2 + \vert B_1^2 \vert^2) \right. \]
\[ + \frac{1}{2} \frac{m_1^2 - m_2^2}{m_{\lambda s}^2} \left[ \sqrt{\hat{r}} (\text{Re}(A_1^* A_2^{* \prime}) + \text{Re}(B_1^* B_2^{* \prime})) - (\text{Re}(A_1^* B_2^{* \prime}) + \text{Re}(B_1^* A_2^{* \prime})) \right] \]
\[ + \frac{1}{2} \frac{m_1^2 - m_2^2}{m_{\lambda s}^2} \left[ \sqrt{\hat{r}} (\text{Re}(A_1^* A_2^{* \prime}) + \text{Re}(B_1^* B_2^{* \prime})) + (\text{Re}(A_1^* B_2^{* \prime}) + \text{Re}(B_1^* A_2^{* \prime})) \right] \]
\[ + \frac{\hat{s}}{2} (m_1^2 - m_2^2)(\text{Re}(A_1^* A_2^{* \prime}) + \text{Re}(B_1^* B_2^{* \prime})) \right], \] (50)
and

\[ I_2(\hat{s}) = \frac{\lambda_2}{\hat{s}^2} \left[ -\frac{1}{4} (|A_1|^2 + |B_1|^2 - m_A^2 \hat{s} (|A_2|^2 + |B_2|^2)) \right]. \quad (51) \]

Here, \( \lambda = \lambda_3 \) (as defined in section 3), \( \lambda_2 = \hat{m}_1^4 + \hat{m}_2^4 + \hat{s}^2 - 2(\hat{m}_1^2 \hat{m}_2^2 + \hat{m}_1^2 \hat{s} + \hat{m}_2^2 \hat{s}) \), and

\[ t = (1 + r - \hat{s})/2. \]

The full kinematically accessible physical range for these processes is given by

\[ (m_1 + m_2)^2 \leq q^2 \leq (m_\Lambda - m_\Lambda)^2. \quad (52) \]

As there is no intermediate particle in the SM which can decay into two leptons of different flavours, so in comparison with the \( \Lambda_b \rightarrow \Lambda^{+}l^{-}\) processes, LFV decays have no long distance QCD contributions and dominant charmonium resonance background. The required input values for numerical evaluation are taken from [8] and the values of the \( q^2 \) dependent form factors are taken from LCSR approach [21]. To determine the values of various LQ couplings, which are involved in the LFV decays, we use the following assumptions. As we know that the expansion parameter of the CKM matrix in the Wolfenstein parametrization (\( \lambda \)), can be related to the down type quark masses as \( \lambda \sim (m_d/m_u)^{1/2} \) in the quark sector, while in the lepton sector one can have the same order for \( \lambda \) with the relation \( \lambda \sim (m_{l_j}/m_{l_i})^{1/2} \). Hence, for other required leptoquark coupling, we assume that the coupling between different generation of quarks and leptons follow the simple scaling laws, i.e., \( \lambda^{(X)} = (m_j/m_i)^{1/2} \lambda^{(X)} \) with \( j > i \). Thus, using the values of the leptoquark coupling as given in table 1, one can obtain the bound on required LQ couplings involved in LFV decays. Using these values we plot the variation of branching ratio of LFV decays such as \( \Lambda_b \rightarrow \Lambda^{+}e^{-} \) (top left panel), \( \Lambda_b \rightarrow \Lambda^{+}\mu^{-} \) (top right panel) and \( \Lambda_b \rightarrow \Lambda^{+}\tau^{-} \) (lower panel) with respect to \( q^2 \) in figure 9 and the predicted upper limits of the branching ratios are given in table 4. So far there is no experimental evidence on the LFV \( \Lambda_b \) decays. However, since the predicted branching ratios are \( O(10^{-9}) \), they can be searched at LHCb and exploration/observation of these modes would definitely shed some light in the leptoquark scenarios.

### 6. Conclusion

In this paper, we have studied the rare semileptonic \( \Lambda_b \rightarrow \Lambda^{+}l^{-} \), \( l = e, \mu, \tau \) baryonic decays in the scalar leptoquark model. The leptoquark parameter space has been constrained using the experimental limits on the branching ratios of the two body leptonic decays \( B_u \rightarrow l^{+}l^{-} \). We have computed the branching ratios, the forward–backward and lepton polarisation asymmetries \( (P_{L,T,N}) \) using the new leptoquark couplings. We have...
shown explicitly the results for both the relevant $X = (3, 2, 7/6)$ and $X = (3, 2, 1/6)$ leptoquark models. The zero-position of the forward–backward asymmetry is found to be insensitive to the additional leptoquark effect. These models also give negligible contribution to the transverse polarisation asymmetry. In addition, we also estimated the LUV parameters in these decays analogous to $R_K$ in $B \to K\pi\pi$ process. The LFV $\Lambda_b$ decays are also studied and the predicted upper limits on these branching ratios are found to be $O(10^{-10} - 10^{-9})$, which could be searched in the LHCb experiment.

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Appendix A. Form factors in $\Lambda_b \to \Lambda$ transition

The transition form factors for $\Lambda_b(P) \to \Lambda(p')$ decays can be parameterised as [16, 43]

$$
\langle \Lambda(p') | \bar{s} \gamma_{\mu} b | \Lambda_b(P) \rangle = f_1 \bar{u}_\Lambda \gamma_{\mu} u_{\Lambda_b} + f_2 \bar{u}_\Lambda \gamma_{\mu} q' u_{\Lambda_b} + f_3 q'_{\mu} A_{\mu b}, \\
\langle \Lambda(p') | \bar{s} \gamma_{\mu} \gamma_5 b | \Lambda_b(P) \rangle = g_1 \bar{u}_\Lambda \gamma_{\mu} \gamma_5 u_{\Lambda_b} + g_2 \bar{u}_\Lambda \gamma_{\mu} q' \gamma_5 u_{\Lambda_b} + g_3 q'_{\mu} A_{\mu b},
$$

and for dipole operators

$$
\langle \Lambda(p') | \bar{s} \gamma_{\mu} q' b | \Lambda_b(P) \rangle = f_1^T \bar{u}_\Lambda \gamma_{\mu} u_{\Lambda_b} + f_2^T \bar{u}_\Lambda \gamma_{\mu} q' u_{\Lambda_b} + f_3^T q'_{\mu} A_{\mu b},
$$

and

$$
\langle \Lambda(p') | \bar{s} \gamma_{\mu} q'' \gamma_5 b | \Lambda_b(P) \rangle = g_1^T \bar{u}_\Lambda \gamma_{\mu} \gamma_5 u_{\Lambda_b} + g_2^T \bar{u}_\Lambda \gamma_{\mu} q'' \gamma_5 u_{\Lambda_b} + g_3^T q'_{\mu} A_{\mu b}.
$$

with $q = P - p'$ and

$$
\begin{align*}
f_2^T &= f_T - f_2^T q^2, \\
f_1^T &= [f_1^V + f_2^S (m_\Lambda + m_{\Lambda_b})] q^2, \\
f_1^T &= - \frac{m_\Lambda^2 - m_{\Lambda_b}^2}{m_\Lambda^2 - m_{\Lambda_b}^2} f_1^T, \\
g_2^T &= g_T - g_2^T q^2, \\
g_1^T &= [g_1^V + g_2^S (m_\Lambda - m_{\Lambda_b})] q^2, \\
g_1^T &= \frac{q^2}{m_\Lambda^2 - m_{\Lambda_b}^2} g_3^T.
\end{align*}
$$

Appendix B. Expressions for $K_{0,1,2}(s)$ functions

The complete expressions for $K_{0,1,2}(s)$ functions required to calculate the double differential decay rate is given by [18]

| Decay process | Predicted branching ratio |
|---------------|--------------------------|
| $\Lambda_b \to \Lambda \mu^- e^+$ | $1.56 \times 10^{-9}$ |
| $\Lambda_b \to \Lambda e^- \mu^+$ | $3.2 \times 10^{-10}$ |
| $\Lambda_b \to \Lambda \mu^- \mu^+$ | $4.6 \times 10^{-9}$ |
\[
\mathcal{K}_0(\bar{s}) = 32 m^2 \bar{s} \left(1 + r - s\right) \left(D_1^2 E_1 + |E_2|^2\right) + 64 m^2 \bar{s} \left(1 - r + s\right) \left(D_2^2 E_1 - D_1^2 E_2 + 6m \bar{s} \sqrt{r} \left(6m^2 - 3m \bar{s}\right) \text{Re}(D_1^2 E_1) + 64m \bar{s} \sqrt{r} \left(2m \bar{s} \Re(D_1^2 E_3) + \left(1 - r + s\right) \left(D_3^2 D_1 + E_4^2 E_3\right)\right) + 32m \bar{s} \sqrt{r} \left(2m \bar{s} \Re(A_2 B_2) - m \bar{s} \Re(A_4 B_2) - 2 \sqrt{r} \left(\Re(A_4 B_3) - m \bar{s} \Re(A_4 B_2)\right)\right) + 8m \bar{s} \sqrt{r} \left(4m^2 \left(1 + r - s\right) \Re(A_4 B_2 + A_4 B_3) - 2 \sqrt{r} \left|\Re(A_4 B_1) + m \bar{s} \Re(A_4 B_2)\right|\right) + 8m \bar{s} \sqrt{r} \left(4m^2 \left(1 + r - s\right) - m \bar{s} \Re(A_4 B_2) - 2 \sqrt{r} \left|\Re(A_4 B_1) + |B_2|^2\right|\right) - 8m \bar{s} \sqrt{r} \left(4m^2 \left(1 + r - s\right) - m \bar{s} \Re(A_4 B_2) - 2 \sqrt{r} \left|\Re(A_4 B_1) + |B_2|^2\right|\right) - 4 \left(1 - r + s\right) \left(D_1^2 E_2 + D_2^2 E_1 + m \bar{s} \Re(A_4 B_2)\right) \left|\left(1 - r + s\right)^2 - \bar{s}^2\right|\left(|D_2|^2 + |E_2|^2\right)\right],
\]

(B.1)

and

\[
\mathcal{K}_2(\bar{s}) = 8 m^4 \bar{s} \left|\left(1 - r + s\right) \left(A_2^2 + |B_2|^2 + |D_2|^2 + |E_2|^2\right) - 8 m^4 \bar{s} \sqrt{r} \left(4m^2 \left(1 + r - s\right) - m \bar{s} \Re(A_4 B_2) - 2 \sqrt{r} \left|\Re(A_4 B_1) + |B_2|^2\right|\right) - 4 \left(1 - r + s\right) \left(D_1^2 E_2 + D_2^2 E_1 + m \bar{s} \Re(A_4 B_2)\right) \left|\left(1 - r + s\right)^2 - \bar{s}^2\right|\left(|D_2|^2 + |E_2|^2\right)\right).\]

(B.3)

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