Lepto-Quark Portal Dark Matter

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Abstract

We consider the extension of the Standard Model with scalar leptoquarks as a portal to dark matter (DM), motivated by the recent anomalies in semi-leptonic $B$-meson decays. Taking singlet and triplet scalar leptoquarks as the best scenarios for explaining $B$-meson anomalies, we discuss the phenomenological constraints from rare meson decays, muon $(g-2)_\mu$, and leptoquark searches at the Large Hadron Collider (LHC). Introducing leptoquark couplings to scalar dark matter, we find that the DM annihilations into a pair of leptoquarks open a wide parameter space, being compatible with XENON1T bound, and show that there is an interesting interplay between LHC leptoquark searches and distinct signatures from cascade annihilations of dark matter.

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1 Introduction

Recently, there have been intriguing anomalies in the semi-leptonic decays of $B$-mesons at BaBar, Belle and LHCb experiments, which are based on the observables of testing Lepton Flavor Universality (LFU), i.e. $R_{K^*(s)}$ [1–3] and $R_{D^*(s)}$ [4–6]. Thus, it is plausible that LFU might be violated due to new physics in the neutral and charged currents associated with muon and tau leptons, respectively. Currently, experimental values of $R_{K^*(s)}$ and $R_{D^*(s)}$ turn out to be deviated from the SM expectations at about 4σ level per each. However, we still need to understand the hadronic uncertainties in angular distributions of related $B$-meson decays [7] and the results are to be confirmed at LHCb with more data and Belle II [8]. Nonetheless, it is important to study the consequences of new physics in direct searches at the LHC and other precision and indirect observables.

Dark matter (DM) is known to occupy about 85% of the total matter density in the Universe, and there are a variety of evidences for the existence of dark matter such as galaxy rotation curves, gravitational lensing, large scale structure, etc. The Weakly Interacting Massive Particles (WIMPs) paradigm has driven forces for searching particle dark matter with non-gravitational interactions beyond the Standard Model (SM) for more than three decades. Various direct detection experiments [9–12] have put stringent bounds on the cross section of DM-nucleon elastic scattering, and forthcoming XENON-nT and large-scale experiments such as DARWIN [13] and LZ [14] will push the limits further to the neutrino floor where there are irreducible backgrounds due to neutrino coherent scattering. In particular, Higgs-portal type models for dark matter have been strongly constrained, apart from the resonance region or the heavy DM masses.

Leptoquark models [15,16] have been revived recently because they can provide an economic way of accommodating the aforementioned $B$-meson anomalies [17–23] and can be tested at the LHC. Leptoquarks carry extra Yukawa-type couplings to the SM fermions, providing a source for violating LFU. Furthermore, leptoquark scalars or vectors could be originated from unified models of forces [24], in analogy to colored triplet Higgs scalars or $X,Y$ gauge bosons in the minimal $SU(5)$ unification. The best scenarios for explaining the $B$-meson anomalies [18,19] are: one $SU(2)_L$-singlet scalar leptoquark $S_1$ for $R_{D^*(s)}$, and one $SU(2)_L$-triplet scalar leptoquark $S_3$ for $R_{K^*(s)}$, or one $SU(2)_L$-singlet vector leptoquark for both $B$-meson anomalies. Leptoquark scenarios are phenomenologically rich, because the muon $(g−2)_\mu$ anomalies can be also explained by leptoquark couplings and various LHC searches can be reinterpreted to bound the leptoquark models.

In this article, we consider a leptoquark-portal model for dark matter where scalar dark matter communicates with the SM through the quartic couplings of scalar leptoquarks, $S_1$ and $S_3$. We show that sizable leptoquark couplings to dark matter lead to new annihilation channels of dark matter into a pair of leptoquarks, opening a wide parameter space where the correct relic density can be explained, being compatible with the direct detection bounds from XENON1T. Moreover, we also discuss that the cascade annihilations of dark matter can lead to distinct signatures for cosmic ray observation, in correlation
to leptoquark searches at the LHC. We argue that our models with scalar leptoquarks are consistent with the current bounds from rare meson decays, mixings and lepton flavor violation, whereas the loop corrections of leptoquarks to DM-nucleon couplings and Higgs couplings can be negligible in most of the parameter space of our interest.

The paper is organized as follows. We first give a brief overview on the $R_{K^{(*)}}$ and $R_{D^{(*)}}$ anomalies and the necessary corrections to the effective Hamiltonians. Then, in models with scalar leptoquarks, we derive the effective interactions for the semi-leptonic $B$-meson decays and discuss the conditions for $B$-meson anomalies and various constraints from rare meson decays, mixings, muon $(g - 2)_\mu$ and leptoquark searches at the LHC. Next we describe leptoquark-portal models for dark matter and consider various constraints on the models, coming from the relic density, direct and indirect detection of dark matter and Higgs data. There are two appendices dealing with the details on effective Hamiltonians for $B$-meson decays and effective interactions for dark matter and Higgs due to leptoquarks, respectively. Finally, conclusions are drawn.

2 Overview on $R_{K^{(*)}}$ and $R_{D^{(*)}}$ anomalies

In this section, we give a brief overview on the status of the $B$-meson anomalies and the interpretations in terms of the effective Hamiltonians in the SM.

The reported value of $R_K = \mathcal{B}(B \to K\mu^+\mu^-)/\mathcal{B}(B \to Ke^+e^-)$ \[1\] is 

$$R_K = 0.745^{+0.090}_{-0.074} \text{(stat)} \pm 0.036 \text{(syst)}, \quad 1 \text{ GeV}^2 < q^2 < 6 \text{ GeV}^2;$$

which deviates from the SM prediction by $2.6\sigma$. On the other hand for vector $B$-mesons, $R_{K^{*}} = \mathcal{B}(B \to K^{*}\mu^+\mu^-)/\mathcal{B}(B \to K^{*}e^+e^-)$ \[2\] is

$$R_{K^{*}} = 0.66^{+0.11}_{-0.07} \text{(stat)} \pm 0.03 \text{(syst)}, \quad 0.045 \text{ GeV}^2 < q^2 < 1.1 \text{ GeV}^2,$$

$$R_{K^{*}} = 0.69^{+0.11}_{-0.07} \text{(stat)} \pm 0.05 \text{(syst)}, \quad 1.1 \text{ GeV}^2 < q^2 < 6.0 \text{ GeV}^2,$$

which again differs from the SM prediction by $2.1–2.3\sigma$ and $2.4–2.5\sigma$, depending on the energy bins. The deviation in $R_{K^{*}}$ is supported by the reduction in the angular distribution of $B \to K^{*}\mu^+\mu^-$, the so called $P'_5$ variable \[3\].

The effective Hamiltonian for $b \to s\mu^+\mu^-$ is given by

$$\Delta H_{\text{eff}, b \to s\mu^+\mu^-} = -\frac{4G_F}{\sqrt{2}} V_{ts}^* V_{tb} \alpha_{em} \frac{\alpha_{em}}{4\pi} (C_9^\mu\mathcal{O}_9^\mu + C_{10}^\mu\mathcal{O}_{10}^\mu + C_{19}^\mu\mathcal{O}_{19}^\mu + C_{10}^\mu\mathcal{O}_{10}^\mu) + \text{h.c.} \quad (3)$$

where $\mathcal{O}_9^\mu \equiv (\bar{s}\gamma^\mu P_L b)(\bar{\mu}\gamma_\mu\gamma^5 \mu)$, $\mathcal{O}_{10}^\mu \equiv (\bar{s}\gamma^\mu P_L b)(\bar{\mu}\gamma_\mu\gamma^5 \mu)$, $\mathcal{O}_9^\mu \equiv (\bar{s}\gamma^\mu P_R b)(\bar{\mu}\gamma_\mu\gamma^5 \mu)$ and $\mathcal{O}_{10}^\mu \equiv (\bar{s}\gamma^\mu P_R b)(\bar{\mu}\gamma_\mu\gamma^5 \mu)$, and $\alpha_{em}$ is the electromagnetic coupling. In the SM, the Wilson coefficients are given by $C_{9,10}^{\mu,SM}(m_b) = -C_{10}^{\mu,SM}(m_b) = 4.27$ and $C_{9,10}^{\mu,SM}(m_b) \approx -C_{10}^{\mu,SM}(m_b) \approx 0.$
For $C_{10}^{\mu,\text{NP}} = C_{9}^{\mu,\text{NP}} = C_{10}^{\mu,\text{NP}} = 0$, the best-fit value for new physics contribution is given by $C_{9}^{\mu,\text{NP}} = -1.11^{25}$, (while taking $[-1.28, -0.94]$ and $[-1.45, -0.75]$ within 1σ and 2σ errors), to explain the $R_{K^{(*)}}$ anomalies. On the other hand, for $C_{9}^{\mu,\text{NP}} = -C_{10}^{\mu,\text{NP}}$ and others being zero, the best-fit value for new physics contribution is given by $C_{9}^{\mu,\text{NP}} = -0.62^{25}$, (while taking $[-0.75, -0.49]$ and $[-0.88, -0.37]$ within 1σ and 2σ errors).

Taking the results of BaBar $^4$, Belle $^5$ and LHCb $^6$ for $R_D = \mathcal{B}(B \to D\tau\nu)/\mathcal{B}(B \to Dl\nu)$ and $R_{D^*} = \mathcal{B}(B \to D^*\tau\nu)/\mathcal{B}(B \to D^*l\nu)$ with $l = e, \mu$ for BaBar and Belle and $l = \mu$ for LHCb, the Heavy Flavor Averaging Group $^{26}$ reported the experimental world averages as follows,

$$R_{D}^{\text{exp}} = 0.403 \pm 0.040 \pm 0.024,$$

$$R_{D^*}^{\text{exp}} = 0.310 \pm 0.015 \pm 0.008.\quad (4)$$

On the other hand, taking into account the lattice calculation of $R_D$, which is $R_D = 0.299 \pm 0.011^{27}$, and the uncertainties in $R_{D^*}$ in various groups $^{28,29}$, we take the SM predictions for these ratios as follows,

$$R_{D}^{\text{SM}} = 0.299 \pm 0.011,$$

$$R_{D^*}^{\text{SM}} = 0.260 \pm 0.010.\quad (6)$$

Then, the combined derivation between the measurements and the SM predictions for $R_D$ and $R_{D^*}$ is about $4.1\sigma$. We quote the best fit values for $R_D$ and $R_{D^*}$ including the new physics contributions $^{30},$

$$\frac{R_D^{\text{SM}}}{R_D^{\text{SM}}} = \frac{R_{D^*}^{\text{SM}}}{R_{D^*}^{\text{SM}}} = 1.21 \pm 0.06.\quad (8)$$

The effective Hamiltonian for $b \to c\tau\nu$ in the SM is given by

$$H_{\text{eff}} = \frac{4G_F}{\sqrt{2}} V_{cb} C_{cb} (\bar{c}\gamma^\mu P_L b)(\bar{\tau}\gamma_\mu P_L \nu_\tau) + \text{h.c.} \quad (9)$$

where $C_{cb} = 1$ in the SM with $V_{cb} \approx 0.04$. The new physics contribution may contain the dimension-6 four-fermion vector operators, $\mathcal{O}_{V_{R,L}} = (\bar{c}\gamma^\mu P_{R,L} b)(\bar{\tau}\gamma_\mu P_L \nu_\tau)$ and/or scalar operators, $\mathcal{O}_{S_{R,L}} = (\bar{c}P_{R,L} b)(\bar{\tau}P_L \nu_\tau)$. Then, in order to explain the $R_{D^{(*)}}$ anomalies in eq. (8), the Wilson coefficient for the new physics contribution should be $\Delta C_{cb} = 0.1$ from eq. (9), while taking $[0.072, 0.127]$ and $[0.044, 0.153]$ within 1σ and 2σ errors.

### 3 Leptoquarks for $B$-meson anomalies

It is known that $SU(2)_L$ singlet and triplet scalar leptoquarks can explain $R_{D^{(*)}}$ and $R_{K^{(*)}}$ anomalies, respectively $^{18,19}$. (See also Ref. $^{17,20,21,23}$.) Thus, in this section, focusing on those scalar leptoquark models, we discuss the phenomenological constraints coming from the $B$-meson anomalies.
3.1 Effective interactions from scalar leptoquarks

We consider the Lagrangian for an $SU(2)_L$ singlet scalar leptoquark $S_1$ with $Y = +\frac{1}{3}$, and an $SU(2)_L$ triplet scalar leptoquark, $S_3 \equiv \Phi_{ab}$ with $Y = +\frac{1}{3}$, as follows,

$$\mathcal{L}_{LQ} = \mathcal{L}_{S_1} + \mathcal{L}_{S_3}$$

(10)

$$\mathcal{L}_{S_1} = -\lambda_{ij} Q^a_{Li} (i\sigma^2)_{ab} S_1 L_{Lj}^b + \text{h.c.}$$

= $$-\lambda_{ij} (Q^C)^a_{Ri} (i\sigma^2)_{ab} S_1 L_{Lj}^b + \text{h.c.}$$

(11)

where $a, b$ are $SU(2)_L$ indices, $\sigma^2$ is the second Pauli matrix and $\psi^C = C\bar{\psi}^T$ is the charge conjugate with $C = i\gamma^0\gamma^2$, and

$$\mathcal{L}_{S_3} = -\kappa_{ij} Q^a_{Li} \Phi_{ab} L_{Lj}^b + \text{h.c.}$$

= $$-\kappa_{ij} (Q^C)^a_{Ri} \Phi_{ab} L_{Lj}^b + \text{h.c.}$$

(12)

with

$$\Phi_{ab} = \begin{pmatrix} \sqrt{2}\phi_3 & -\phi_2 \\ -\phi_2 & -\sqrt{2}\phi_1 \end{pmatrix}$$

(13)

where $(\phi_1, \phi_2, \phi_3)$ forms an isospin triplet with $T_3 = +1, 0, -1$ and $Q = +\frac{4}{3}, +\frac{1}{3}, -\frac{2}{3}$. We note that our conventions are comparable to those in the literature by writing $\Phi = (i\sigma^2)(\bar{\sigma} \cdot \bar{S})$ where $\bar{\sigma}$ are Pauli matrices and $\bar{S}$ are complex scalar fields.

Then, after integrating out the leptoquark scalars, we obtain the effective Lagrangian for the SM fermions in the following,

$$\mathcal{L}_{\text{eff}} = \left( \frac{1}{4m_{S_1}^2} \lambda_{ij} \lambda_{kl}^* \right) \left( \bar{Q}_{Li} \gamma^\mu Q_{Li} \right) \left( \bar{L}_{Li} \gamma_{\mu} L_{Lj} \right)$$

$$+ \left( -\frac{1}{4m_{S_3}^2} \lambda_{ij} \lambda_{kl}^* \right) \left( \bar{Q}_{Li} \gamma^\mu \sigma^I Q_{Li} \right) \left( \bar{L}_{Li} \gamma_{\mu} \sigma^I L_{Lj} \right)$$

(14)

where $\sigma^I (I = 1, 2, 3)$ are the Pauli matrices. There, we find that there are both $SU(2)_L$ singlet and triplet $V-A$ operators. As compared to the case with $U(2)$ flavor symmetry [19], the effective interactions for either singlet or triplet leptoquark can be written as

$$\mathcal{L}_{\text{eff}} = -\frac{1}{v^2} \lambda_{ab}^2 \lambda_{ij}^2 \left[ C_S \left( \bar{Q}_{Li} \gamma^\mu Q_{Li} \right) \left( \bar{L}_{Li} \gamma_{\mu} L_{Lj} \right) + C_T \left( \bar{Q}_{Li} \gamma^\mu \sigma^I Q_{Li} \right) \left( \bar{L}_{Li} \gamma_{\mu} \sigma^I L_{Lj} \right) \right]$$

(15)

So, we obtain $C_S = -C_T$ for the singlet leptoquark and $C_S = 3C_T$ for the triplet leptoquark. A fit to low-energy data including the $R_{K^(*)}$ and $R_{D^(*)}$ anomalies has been done with four free parameters, $C_T, C_S, \lambda_{ab}^2$ and $\lambda_{ij}^2$, under the assumption that the CKM matrix stems solely from the mixing between up-type quarks [19]. As a result, the best-fit values are given by $C_S \approx C_T \approx 0.02$ for $|\lambda_{ab}^2| < 5|V_{ub}|$ [19].
Figure 1: Parameter space for the leptoquark mass $m_{LQ}$ and the effective coupling $\lambda_{\text{eff}}$, explaining the B-meson anomalies, in green(yellow) region at 2σ(1σ) level. We have taken $m_{LQ} = m_{S_1}$ and $\lambda_{\text{eff}} = \sqrt{|\lambda_{33}^{*}\lambda_{23}|}$ for $R_{D(\ast)}$ on left plot, and $m_{LQ} = m_{S_3}$ and $\lambda_{\text{eff}} = \sqrt{|\kappa_{32}^{*}\kappa_{22}|}$ for $R_{K(\ast)}$ on right plot.

### 3.2 Singlet scalar leptoquark

After integrating out the leptoquark $S_1$, from the results in eq. (A.2), we obtain the effective Hamiltonian relevant for $b \rightarrow c\tau\bar{\nu}_\tau$ as

$$H_{S_1}^{b \rightarrow c\tau\bar{\nu}_\tau} = \frac{\lambda_{33}^{*}\lambda_{23}}{2m_{S_1}^2} (\bar{b}_L\gamma^\mu c_L)(\bar{\nu}_{\tau L}\gamma^\mu \tau_L) + \text{h.c.} \equiv \frac{1}{\Lambda_D^2} (\bar{b}_L\gamma^\mu c_L)(\bar{\nu}_{\tau L}\gamma^\mu \tau_L) + \text{h.c.}.$$  

(16)

As a consequence, the singlet leptoquark gives rise to the effective operator for explaining the $R_{D(\ast)}$ anomalies and and the effective cutoff scale is to be $\Lambda_D \sim 3.5 \text{ TeV}$ [31]. Thus, for $m_{S_1} \gtrsim 1 \text{ TeV}$, we need $\sqrt{|\lambda_{33}^{*}\lambda_{23}|} \gtrsim 0.4$.

In the left plot of Fig. 1, we depict the parameter space for $m_{S_1}$ and the effective leptoquark coupling, $\lambda_{\text{eff}} = \sqrt{|\lambda_{33}^{*}\lambda_{23}|}$, in which the $R_{D(\ast)}$ anomalies can be explained within 2σ(1σ) errors in green(yellow) region from the conditions below eq. (9).

From the couplings of the singlet scalar leptoquark necessary for $R_{D(\ast)}$ anomalies,

$$\mathcal{L}_{S_1} \supset -\lambda_{33} \left( \overline{(t^C)R} S_1 \tau_L - \overline{(b^C)R} S_1 \nu_{\tau L} \right) + \text{h.c.}$$

$$-\lambda_{23} \left( \overline{(c^C)R} S_1 \tau_L - \overline{(s^C)R} S_1 \nu_{\tau L} \right) + \text{h.c.},$$

(17)

the decay modes of the singlet scalar leptoquark are given by $S_1 \rightarrow \overline{t}\tau, \overline{c}\tau$ and $S_1 \rightarrow \overline{b}\nu_{\tau}, \overline{s}\nu_{\tau}$, which are summarized together with the corresponding LHC bounds on leptoquark masses in Table II.
The decay branching ratios of the triplet leptoquark and the corresponding $\phi$ and $\gamma$ are also obtained. The triplet leptoquark gives rise to the effective operator of the $(V-A)$ form for the quark current, that is, $\mathcal{H}_b \rightarrow s\mu^+\mu^-$ as

$$\mathcal{H}_{b \rightarrow s\mu^+\mu^-} = -\frac{\kappa_{22}^2}{m_{\phi_1}^2} \left( \bar{b}_L \gamma^\mu s_L \right) \left( \bar{\mu}_L \gamma_\mu \mu_L \right) + \text{h.c.} \equiv \frac{1}{\Lambda_K^2} \left( \bar{b}_L \gamma^\mu s_L \right) \left( \bar{\mu}_L \gamma_\mu \mu_L \right) + \text{h.c.} \quad (18)$$

As a consequence, the triplet leptoquark gives rise to the effective operator of the $(V - A)$ form for the quark current, that is, $C_{9,10}^{\mu,\text{NP}} = -C_{9,10}^{\mu,\text{NP}} \neq 0$, as favored by the $R_{K^{(*)}}$ anomalies, and the effective cutoff scale is to be $\Lambda_K \sim 30 \text{TeV}$ [37]. The result is in contrast to the case for $Z'$ models with family-dependent charges such as $Q' = x(B_3 - L_3) + y(L_\mu - L_\tau)$ with $x, y$ being arbitrary parameters where $C_9^{\mu,\text{NP}} \neq 0$ and $C_{10}^{\mu,\text{NP}} = 0$ [39]. Then, for $m_{\phi_1} \gtrsim 1 \text{TeV}$, we need $\sqrt{\kappa_{22}^2} \gtrsim 0.03$. Therefore, we can combine scalar leptoquarks, $S_1$ and $S_3$, to explain $R_{D^{(*)}}$ and $R_{K^{(*)}}$ anomalies, respectively.

In the right plot of Fig. 1, we depict the parameter space for $m_{S_1}$ and the effective leptoquark coupling, $\lambda_{\text{eff}} = \sqrt{\kappa_{22}^2}$, in which the $R_{K^{(*)}}$ anomalies can be explained within $2\sigma (1\sigma)$ errors in green(yellow) region from the conditions below eq. (3).

Likewise as for the singlet scalar leptoquark, from the triplet leptoquark couplings necessary for $R_{K^{(*)}}$ anomalies,

$$\mathcal{L}_{S_3} \supset -\kappa_{32} \left( \sqrt{2} (t^C)_R \phi_3 \nu_{\mu L} - (t^C)_R \phi_2 \mu_L - (b^C)_R \phi_2 \nu_{\mu L} - \sqrt{2} (b^C)_R \phi_1 \mu_L \right) + \text{h.c.} \quad (19)$$

the decay modes of the singlet scalar leptoquark are given by $\phi_1 \rightarrow \bar{b}\mu, \bar{s}\mu, \phi_2 \rightarrow \bar{t}\mu, \bar{c}\mu, \bar{b}\nu_\mu, \bar{s}\nu_\mu, \text{ and } \phi_3 \rightarrow \bar{\tau}\nu_\mu, \bar{\nu}_\tau$. As will be discussed in the next section, the bounds from $B \rightarrow K\nu\bar{\nu}$ could require $\kappa_{33}$ and $\kappa_{23}$ to be sizable. In this case, the decay modes containing $\bar{\tau}$ or $\nu_\tau$ are relevant too. The decay branching ratios of the triplet leptoquark and the corresponding LHC bounds on the mass of triplet scalar leptoquark are also summarized in Table I.

| LQs | BRs | $m_{LQ,\text{min}}$ | BRs | $m_{LQ,\text{min}}$ |
|-----|-----|---------------------|-----|---------------------|
| $S_1$ | $B(\bar{t}\tau/\nu_{\tau}) = \frac{1}{2} \beta$ | 1.22 TeV [32] | $B(\bar{c}\tau/\nu_{\tau}) = \frac{1}{2} (1 - \beta)$ | 950 GeV ($\nu_{\tau} J$) [33] |
| $S_3(\phi_1)$ | $B(\bar{b}\mu) = \gamma$ | 1.4 TeV [34] | $B(s\mu) = 1 - \gamma$ | 1.08 TeV ($\mu J$) [35] |
| $S_3(\phi_2)$ | $B(t\mu/\nu_\mu) = \frac{1}{2} \gamma$ | 1.45 TeV ($t\mu$) [36] | $B(\bar{c}\mu/\nu_\mu) = \frac{1}{2} (1 - \gamma)$ | 850 GeV ($\mu\nu_\mu JJ$) [37] |
| $S_3(\phi_3)$ | $B(t\nu_\mu) = \gamma$ | 1.12 TeV [38] | $B(\bar{c}\nu_\mu) = 1 - \gamma$ | 950 GeV ($\nu_\mu JJ$) [38] |

Table 1: Decay branching ratios of leptoquarks, and LHC bounds on leptoquark masses. Here, $\beta \equiv \lambda_{33}^2 / (\lambda_{23}^2 + \lambda_{23}^2)$ and $\gamma \equiv \kappa_{22}^2 / (\kappa_{22}^2 + \kappa_{23}^2)$. Most LHC bounds are given for $B = 1$, except in Ref. [37] where $B(\bar{c}\mu) = B(s\nu_\mu) = 0.5$ was taken.
4 Constraints on leptoquarks

We discuss the constraints on scalar leptoquark models, due to other rare meson decays, muon \((g-2)_\mu\), lepton flavor violation as well as the LHC searches.

4.1 Rare meson decays and mixing

In leptoquark models explaining the B-meson anomalies, there is no \(B - \bar{B}\) mixing at tree level, but instead it appears at one-loop level. Therefore, the resulting new contribution to the \(B_s - \bar{B}_s\) mixing is about 1% level \([17]\), which can be ignored.

Both singlet and triplet leptoquarks contribute to \(B \rightarrow K(\ast)\nu\bar{\nu}\) at tree level, so their couplings are severely constrained in this case \([17,19]\). The effective Hamiltonian relevant for \(\bar{b} \rightarrow \bar{s}\nu\bar{\nu}\)\([40]\) is

\[
\mathcal{H}_{\bar{b} \rightarrow \bar{s}\nu\bar{\nu}} = -\frac{\sqrt{2}\alpha_{em}G_F}{\pi} V_{tb}V_{ts}^* \sum_l C_l^d (\bar{b}\gamma^\mu P_L s)(\bar{\nu}_l\gamma_\mu P_L \nu_l)
\]

where \(C_L^d = C_{L,SM} + C_{L,\text{NP}}\). Here, the SM contribution \(C_{L,SM}\) is given by

\[
C_{L,SM} = -\frac{X_t}{s_W^2} \frac{\pi}{\sqrt{2}\alpha_{em}G_FV_{tb}V_{ts}^*}
\]

where \(s_W \equiv \sin \theta_W\) and \(X_t = 1.469 \pm 0.017\). From the result in eq. (A.9), the scalar leptoquarks leads to additional contributions to the effective Hamiltonian for \(B \rightarrow K\nu\bar{\nu}\) as

\[
C_{\nu}^{NP} = -\left(\frac{\lambda_3^*\lambda_{2j}}{2m_{S_1}^2} + \frac{\kappa_{3j}\kappa_{2j}}{2m_{\phi_2}^2}\right)\frac{\pi}{\sqrt{2}\alpha_{em}G_FV_{tb}V_{ts}^*}
\]

Therefore, the ratio of the branching ratios are given by

\[
R_{K(\ast)\nu} \equiv \frac{B(B \rightarrow K(\ast)\nu\bar{\nu})}{B(B \rightarrow K(\ast)\nu\bar{\nu})_{SM}} = \frac{2}{3} + \frac{1}{3} \frac{|C_{L,SM} + C_{\nu,\text{NP}}|^2}{|C_{L,SM}|^2}
\]

Comparing the experimental bounds on \(B(B \rightarrow K(\ast)\nu\bar{\nu})\) \([41]\) given by

\[
B(B \rightarrow K\nu\bar{\nu}) < 1.6 \times 10^{-5}, \quad B(B \rightarrow K^*\nu\bar{\nu}) < 2.7 \times 10^{-5}
\]

to the SM values \([42]\) given by

\[
B(B \rightarrow K\nu\bar{\nu})_{SM} = (3.98 \pm 0.43 \pm 0.19) \times 10^{-6}, \quad B(B \rightarrow K^*\nu\bar{\nu})_{SM} = (9.19 \pm 0.86 \pm 0.50) \times 10^{-6}
\]
and ignoring the imaginary part of $C^t_{\nu}^{NP}$, we get the $R_{K^*\nu}$ bound as

$$-10.1 < \text{Re}(C^t_{\nu}^{NP}) < 22.8.$$  \hfill (25)

Taking into account $\kappa_{32}$ and $\kappa_{23}$, which are necessary for $B \to K^{(*)}\mu^+\mu^-$, the triplet scalar leptoquark contributes only to $B \to K^{(*)}\nu_{\mu}\bar{\nu}_{\mu}$. In this case, as the triplet leptoquark contribution to $C^\nu_{\nu}^{NP}$ is about the same as $C^\nu_{\nu}^{NP} = -0.61$, it satisfies the $R_{K^*\nu}$ bound on its own easily.

On the other hand, the singlet leptoquark with nonzero $\lambda_{33}$ and $\lambda_{23}$, which are necessary for $B \to D^{(*)}\tau\bar{\nu}_\tau$, contribute significantly to $B \to K^{(*)}\nu_{\tau}\bar{\nu}_\tau$. Therefore, we need to cancel the singlet scalar leptoquark contributions to $B \to K^{(*)}\nu_{\mu}\bar{\nu}_{\mu}$, by imposing that

$$|\kappa_{33}^*\kappa_{23}| \approx |\kappa_{33}^*\kappa_{23}| \frac{m^2_{S_3}}{m^2_{S_1}} \approx \frac{|\kappa_{33}^*\kappa_{23}|}{m^2_{S_1}}.$$  \hfill (26)

Ignoring the mass splitting generated within the triplet scalar leptoquark due to potential higher dimensional operators after electroweak symmetry breaking, we get $m_{S_1} = m_{S_2} = m_{S_3} = m_{S_1}$. Then, in order to cancel the contributions to $B \to K^{(*)}\nu_{\mu}\bar{\nu}_{\mu}$ or $B \to K^{(*)}\nu_{\mu}\bar{\nu}_{\mu}$, the necessary conditions for the additional couplings are

$$|\kappa_{33}^*\kappa_{23}| \approx \frac{2m^2_{S_3}}{\Lambda^2_D} \lesssim 1,$$  \hfill (27)

$$|\lambda_{32}^*\lambda_{23}| \approx |\lambda_{32}^*\lambda_{23}| \left(\frac{m^2_{S_1}}{m^2_{S_3}}\right).$$  \hfill (28)

Therefore, for $|\kappa_{33}^*\kappa_{23}| = \mathcal{O}(1)$ and $m_{S_3} \sim m_{S_1}$, the additional coupling, $|\lambda_{32}^*\lambda_{23}|$, from the singlet scalar leptoquark, can be sizable.

4.2 $(g - 2)\mu$

For the singlet scalar leptoquark, the relevant Yukawa couplings for $(g - 2)\mu$ with an additional Yukawa coupling, are given as follows,

$$\mathcal{L}_{S_1} \supset -\lambda_{ij}(Q^C)^a_i\Re\text{(i}\sigma^2)_{ab}S_1^a_{Li}\lambda_{ij}(u^C)_{Li}S_1^a_{jR} + h.c.$$  \hfill (29)

Then, the chirality-enhanced effect from the top quark contributes most \cite{17}, as follows,

$$a^S_{\mu} \approx \frac{m_{\mu}}{4\pi^2} \text{Re}[C^{22}_{R}]$$  \hfill (30)

with

$$C^{ij}_{R} \equiv -\frac{N_c}{12m^2_{S_1}} m_t\lambda_{3i}\lambda_{3j}^* \left(7 + 4\log\left(\frac{m^2_{t}}{m^2_{S_1}}\right)\right).$$  \hfill (31)
Figure 2: Parameter space for $m_{LQ} = m_{S_1}$ and $\lambda'_{32}$ allowed by $(g - 2)_\mu$, in green(yellow) region, at $2\sigma(1\sigma)$ level. The gray region is excluded by the bound on $\text{BR}(\tau \to \mu\gamma)$. We have fixed $\lambda_{32} = \lambda_{33} = 1(0.1)$ on left(right) plot and $\lambda'_{33} = 0$ in both plots.

The deviation of the anomalous magnetic moment of muon between experiment and SM values is given \([44,45]\) by

$$\Delta a_\mu = a_{\text{exp}} - a_{\text{SM}} = 288(80) \times 10^{-11},$$

which is a $3.6\sigma$ discrepancy from the SM \([45]\). We note that as discussed in eq. (28), the extra coupling for the triplet leptoquark, $\kappa_{23}$, allows for a sizable $\lambda_{32}$, leading to a large deviation in $(g - 2)_\mu$ without a conflict to the bound from $B(B \to K^{(*)}\nu\bar{\nu})$.

On the other hand, the additional coupling also contributes to the branching ratio of $\tau \to \mu\gamma$ as follows,

$$\text{BR}(\tau \to \mu\gamma) = \frac{\alpha m_\tau^3}{256\pi^4} \tau_\tau \left(|C_{R}^{23}|^2 + |C_{L}^{23}|^2\right)$$

where $C_{L}^{ij} = C_{R}^{ij}(\lambda_{3i} \to \lambda'_{3i}, \lambda'_{3j} \to \lambda_{3j})$ and the lifetime of tau is given by $\tau_\tau = (290.3 \pm 0.5) \times 10^{-15}$ s \([45]\). The current experimental bound is given \([46]\) by

$$\text{BR}(\tau \to \mu\gamma) < 4.4 \times 10^{-8}.$$
4.3 Leptoquark searches

There are two main production channels for leptoquarks at the LHC, one is pair production via gluon fusion and the other is single production via gluon-quark fusion [16,43].

In the case of $R_{K^{(*)}}$ anomalies, the triplet scalar leptoquark ($\phi_1$) couples to $b/s, \mu$. The other components of the triplet leptoquark couple to $b/s, \nu_\tau$ and $t/c, \mu$ for $\phi_2$ and $t/c, \nu_\mu$ for $\phi_3$. On the other hand, in the case of $R_{D^{(*)}}$ anomalies, the singlet scalar leptoquark ($S_1$) couples to $b/s, \nu_\tau$ and $t/c, \tau$. When the leptoquark pair production via gluon fusion is dominant, the current limits on leptoquark masses listed in Table 1 apply. The current LHC bounds on leptoquarks depend on decay modes, but the leptoquark masses are constrained to be greater than about 1 TeV in most cases.

When the Yukawa couplings, $\phi_1$-$b/s$-$\mu$, $S_1$-$b$-$\nu_\tau$ and $S_1$-$c$-$\tau$ couplings, present in models explaining the $B$-anomalies, are sizable, the leptoquarks can be singly produced by $b/s/c$ quark fusions with gluons. For instance, in the case of $\phi_1$, the relevant production/decay channels are $pp \rightarrow \phi_1^*\phi_1 = bb(s)\mu^+\mu^-$ and $pp \rightarrow \phi_1^+ \rightarrow b(s)\mu^+\mu^- [16]$.

5 Leptoquarks and scalar dark matter

We introduce a scalar dark matter that have direct interactions to scalar leptoquarks and the SM Higgs doublet $H$ by quartic couplings. Thus, this is the minimal dark matter model without a need of extra mediator particle. In this section, we regard scalar leptoquarks as portals to scalar dark matter and discuss the impacts of leptoquarks on direct and indirect detection of dark matter as well as Higgs data.

We can also consider leptoquark-portal models for fermion or vector dark matter too. But, in this case, there is a need of mediator particles [47] and/or non-renormalizable interactions [48], leading to more parameters in the model, so this case is postponed to a future publication for comparison [49].

5.1 Annihilation cross sections for scalar dark matter

We consider a scalar leptoquark $S_{LQ} = S_{1,3}$ and a singlet real scalar dark matter $S$. Then, the most general renormalizable Lagrangian consistent with $S \rightarrow -S$ is

$$L_S = |D_\mu S_{LQ}|^2 - m_{LQ}^2 |S_{LQ}|^2 + \frac{1}{2} (\partial_\mu S)^2 - \frac{1}{2} m_S S^2$$

$$-\frac{1}{4} \lambda_1 S^4 - \frac{1}{2} \lambda_2 |S_{LQ}|^4 - \frac{1}{2} \lambda_3 S^2 |S_{LQ}|^2 - \frac{1}{2} \lambda_4 S^2 |H|^2 - \lambda_5 |H|^2 |S_{LQ}|^2. \quad (35)$$

The above Lagrangian generalizes the Higgs portal interactions to those for leptoquarks.
Figure 3: Feynman diagrams for annihilations of scalar dark matter at tree level.

After electroweak symmetry breaking with $H = (0, v + h)^T / \sqrt{2}$, the new interactions relevant for $SS \rightarrow S_{LQ} S_{LQ}^* h h$ are

$$\mathcal{L}_{S,\text{int}} = -\frac{1}{2} \lambda_3 S^2 |S_{LQ}|^2 - \frac{1}{4} \lambda_4 S^2 (2vh + h^2) - \frac{1}{2} \lambda_5 |S_{LQ}|^2 (2vh + h^2).$$

(36)

Scalar dark matter $S$ annihilates through three channels at tree level, $SS \rightarrow f\bar{f}$ with $f$ being the SM fermions, $SS \rightarrow hh$, with $h$ being the SM Higgs boson, $SS \rightarrow VV$ with $V$ being electroweak gauge bosons, and $SS \rightarrow S_{LQ} S_{LQ}^*$ for $m_{LQ} < m_S$. Depending on the quartic couplings, a heavy scalar dark matter may annihilate into a pair of leptoquarks dominantly, leaving the signatures in both anti-proton and positron from cosmic rays, due to the decay products of leptoquarks, as will be discussed later.

For $m_{LQ} > m_S$, $SS \rightarrow S_{LQ} S_{LQ}^*$ channels are kinematically closed, so instead leptoquark loops make corrections to $SS \rightarrow VV$ with $V$ being electroweak gauge bosons and contribute to new annihilations such as $SS \rightarrow gg, Z\gamma, \gamma\gamma$. In this case, depending on the relative contributions of $SS \rightarrow f\bar{f}, hh, VV$ channels, the loop-induced annihilation channels can be relevant.

We obtain the effective interactions between scalar dark matter and the SM gauge bosons due to leptoquarks with $m_{LQ} > m_S$, as follows,

$$\mathcal{L}_{S,\text{eff}} = D_3 S^2 G_{\mu\nu} G^{\mu\nu} + D_2 S^2 W_{\mu\nu} W^{\mu\nu} + D_1 S^2 F_{Y\mu\nu} F^{Y\mu\nu}$$

(37)

The details on the above effective interactions are given in Appendix B. Then, in the basis of mass eigenstates, the above effective interactions become

$$\mathcal{L}_{S,\text{eff}} = D_{gg} S^2 G_{\mu\nu} G^{\mu\nu} + D_{WW} S^2 W_{\mu\nu}^+ W^{-\mu\nu} + D_{ZZ} S^2 Z_{\mu\nu} Z^{\mu\nu} + D_{\gamma\gamma} S^2 F_{\mu\nu} F^{\mu\nu}$$

(38)
where

\[ D_{gg} = D_3, \]
\[ D_{WW} = 2D_2, \]
\[ D_{ZZ} = D_1 \sin^2 \theta_W + D_2 \cos^2 \theta_W, \]
\[ D_{Z\gamma} = (D_2 - D_1) \sin(2\theta_W), \]
\[ D_{\gamma\gamma} = D_1 \cos^2 \theta_W + D_2 \sin^2 \theta_W. \]

First, the tree-level annihilation cross sections are

\[ (\sigma_{v\text{rel}})_{SS \rightarrow S_L Q S_L^* L_Q} = \frac{N_c N_{LQ}}{32 \pi m_S^2} \left( 1 - \frac{m_{LQ}^2}{m_S^2} \right)^2 \left( \lambda_3 + \frac{\lambda_4 \lambda_5 v^2}{4m_S^2 - m_h^2} \right)^2, \]  
\[ (\sigma_{v\text{rel}})_{SS \rightarrow h h} = \frac{\lambda_4^2}{64 \pi m_S^2} \left( 1 - \frac{m_h^2}{m_S^2} \right)^2 \left( 1 + \frac{3m_h^2}{4m_S^2 - m_h^2} - \frac{2\lambda_4 v^2}{2m_h^2 - m_S^2} \right)^2, \]  
\[ (\sigma_{v\text{rel}})_{SS \rightarrow f \bar{f}} = \frac{N_c \lambda_4^2}{4\pi} \frac{m_f^2}{(4m_S^2 - m_f^2)^2} \left( 1 - \frac{m_f^2}{m_S^2} \right)^{3/2}, \]

with \( f \) being all the SM fermions satisfying \( m_f < m_S \). Here, we note that \( N_c = 3 \) is the number of colors and \( N_{LQ} = 1, 3 \) for \( S_{LQ} = S_1, S_3 \), respectively.

On the other hand, for \( m_{LQ} > m_S \), instead of \( SS \rightarrow S_{LQ} S_{LQ}^* \), we need to consider the loop-induced annihilation cross sections \[50,51\] for \( SS \rightarrow gg, \gamma\gamma, Z\gamma, \) as follows,

\[ (\sigma_{v\text{rel}})_{SS \rightarrow gg} = \frac{64 D_{gg} m_S^2}{\pi}, \]
\[ (\sigma_{v\text{rel}})_{SS \rightarrow \gamma\gamma} = \frac{8 D_{\gamma\gamma}^2 m_S^2}{\pi}, \]
\[ (\sigma_{v\text{rel}})_{SS \rightarrow Z\gamma} = \frac{4 D_{Z\gamma}^2 m_S^2}{\pi} \left( 1 - \frac{m_Z^2}{4m_S^2} \right)^3. \]

Adding loop corrections of leptoquarks to tree-level contributions coming from the Higgs portal coupling \( \lambda_4 \), we also obtain the annihilation cross sections for \( SS \rightarrow WW, ZZ \), respectively,

\[ (\sigma_{v\text{rel}})_{SS \rightarrow WW} = \left\lfloor \frac{\lambda_4^2 m_S^2}{2\pi (m_h^2 - 4m_S^2)^2} \left( 1 - \frac{m_W^2}{m_S^2} + \frac{3m_W^4}{4m_S^4} \right) + \frac{4|D_{WW}|^2 m_S^2}{\pi} \left( 1 - \frac{m_W^2}{m_S^2} + \frac{3m_W^4}{8m_S^4} \right) \right\rfloor \sqrt{1 - \frac{m_W^2}{m_S^2}} \]
\[ + \frac{3\lambda_4 \text{Re}[D_{WW}] m_W^2}{2\pi (m_h^2 - 4m_S^2)} \left( 2 - \frac{m_W^2}{m_S^2} \right) \right\rfloor \sqrt{1 - \frac{m_W^2}{m_S^2}}, \]

(50)
Figure 4: Branching ratios of annihilation cross sections for dark matter as a function of $\lambda_4$ (upper panel) or $m_S$ (lower panel), in models with singlet scalar leptoquark. Branching ratios for $WW$ (orange), $ZZ$ (purple), $gg$ (green), $hh$ (black), $f \bar{f}$ (blue), and $S_{LQ} S^{*}_{LQ}$ (red) channels are shown.

$$
(\sigma v_{\text{rel}})_{SS \rightarrow ZZ} = \left[ \frac{\lambda_4^2 m_S^2}{4\pi(m_h^2 - 4m_S^2)^2} \left( 1 - \frac{m_Z^2}{m_S^2} + \frac{3m_Z^4}{4m_S^4} \right) + \frac{8|D_{ZZ}|^2 m_S^2}{\pi} \left( 1 - \frac{m_Z^2}{m_S^2} + \frac{3m_Z^4}{8m_S^4} \right) \right] \sqrt{1 - \frac{m_Z^2}{m_S^2}}.
$$

(51)

In Figs. 4 and 5, we show the branching ratios of annihilation cross sections of dark matter, $\text{BR}(SS \rightarrow ij)$, as a function of $\lambda_4$ in the upper panel and $m_S$ in the lower panel, for singlet and triplet scalar leptoquarks, respectively.

For light dark matter with $m_S < m_{LQ}$, we find that the tree-level annihilation processes such as $WW, hh, f \bar{f}, ZZ$ are dominant and the loop-induced processes due to leptoquarks are suppressed, except the $gg$ channel, which can be as large as $1 - 10\%$ of the total
Figure 5: The same as in Fig. 4 but in models with triplet scalar leptoquark.
annihilation cross section, depending on whether the scalar leptoquark is singlet or triplet. In the case of triplet scalar leptoquark, the $Z\gamma, \gamma\gamma$ channels can be as large as 1% or 0.1% of the total annihilation cross section, so they could be probed by Fermi-LAT [53] or HESS [54] line searches. On the other hand, for heavy dark matter with $m_S > m_{LQ}$, the $S_{LQ}S_{LQ}^*$ channel becomes dominant while the other tree-level processes are negligible as far as $|\lambda_3| \gtrsim |\lambda_4|$.

5.2 Direct detection bounds

For scalar dark matter, the effective DM-quark interaction is induced due to the SM Higgs exchange at tree level, as follows,

$$L_{\text{eff}, Sq} = \frac{\lambda_4 m_q}{m_h^2} S^2 \bar{q}q.$$  (52)

Moreover, taking a small momentum transfer for the DM-nucleon scattering in eq. (B.2), the effective interactions between scalar dark matter and gluons, generated by loop corrections with leptoquarks, become

$$L_{\text{eff}, Sgg} = \frac{\alpha_s \lambda_4}{96\pi m_{LQ}^2} l_3(S_{LQ}) S^2 G_{\mu\nu}G^{\mu\nu}.$$  (53)

where $l_3(S_{LQ})$ is the Dynkin index of a leptoquark $S_{LQ}$ under $SU(3)_C$. Then, the spin-independent cross section for DM-nucleon elastic scattering is given by

$$\sigma_{S-N} = \frac{\mu_N^2}{\pi m_{LQ}^2} \left(Z f_p + (A - Z) f_n\right)^2.$$  (54)

where $Z, A - Z$ are the numbers of protons and neutrons in the detector nucleus, $\mu_N = m_N m_S/(m_N + m_S)$ is the reduced mass of DM-nucleon system, and

$$f_{p,n} = \frac{\lambda_4 m_{p,n}}{m_h^2} \left(\sum_{q=u,d,s} f_{pq}^n + \frac{2}{9} f_{TG}^n\right) - \frac{\lambda_3 m_{p,n}}{108 m_{LQ}^2} l_3(S_{LQ}) f_{TG}^n.$$  (55)

with $f_{TG}^n = 1 - \sum_{q=u,d,s} f_{Tq}^n$. Here, the mass fractions are $f_{Tu}^p = 0.023, f_{Td}^p = 0.032$ and $f_{Tu}^n = 0.020$ for a proton and $f_{Td}^n = 0.017, f_{Tu}^n = 0.041$ and $f_{Td}^n = 0.020$ for a neutron [52]. Therefore, the quartic coupling $\lambda_4$ between scalar dark matter and SM Higgs is strongly constrained by direct detection experiments such as XENON1T [9]. Consequently, tree-level annihilations of scalar dark matter into $hh, f\bar{f}, WW, ZZ$ are constrained, while the leptoquark-induced annihilations at tree or loop levels can be relevant.

In Fig. 6, we show the DM relic density as a function of DM mass in red solid (dashed) lines for triplet (singlet) scalar leptoquarks. We also show the DM-nucleon scattering cross section in blue lines as can be read from the right vertical axis, and the XENON1T bound in purple dot-dashed lines. We find that the extra annihilation of dark matter into a pair of leptoquarks opens a new parameter space at $m_S > m_{LQ}$ due to a sizable leptoquark portal coupling, $\lambda_3$, avoiding the direct detection bound from XENON1T.
Figure 6: Dark matter relic density as a function of $m_S$ in red solid (dashed) lines for triplet (singlet) scalar leptoquarks. DM-nucleon scattering cross section and XENON1T bound are shown in blue line and purple dot-dashed line, respectively. $\lambda_3 = 0.01, 0.1, 1$ are taken from the top left plot clockwise, and $\lambda_4 = 0.1, \lambda_5 = 1$ and $m_{LQ} = 1$ TeV are taken for all plots.
5.3 Indirect detection bounds

For relatively light scalar dark matter with \( m_S < m_{LQ} \), the DM annihilation cross sections into \( hh, WW, ZZ, tt, bb \) are dominant. In this case, Fermi-LAT dwarf galaxies \(^{55}\) and HESS gamma-rays \(^{56}\) and AMS-02 antiprotons \(^{57}\) can constrain the model.

In Fig. 7, we depict the parameter space in \( \lambda_4 \) vs \( m_S \) in black and red lines, satisfying the correct relic density for models with singlet and triplet scalar leptoquarks, respectively. In the same plots, we superimpose the indirect detection bounds from Fermi-LAT and HESS gamma-ray searches and AMS-02 anti-proton as well as the direct detection bounds from XENON1T. Moreover, the region with \( m_S < m_h/2 \) can be also constrained by Higgs data such as Higgs invisible decay and the signal strength for \( gg \to h \to \gamma\gamma \), as will be discussed in the next subsection.

As a result, the Higgs data as well as indirect detection constrains the region with light and weak-scale dark matter, but the XENON1T experiment constrains most, ruling out most of the DM masses below \( m_S = 1 \) TeV, except the resonance region near \( m_S = m_h/2 \). However, we find that the correct relic density can be obtained for a small value of \( \lambda_4 \) due to the contribution of DM annihilation channels into a leptoquark pair with a sizable leptoquark-portal coupling \( \lambda_3 \) for \( m_S > m_{LQ} \). Therefore, there is a wide parameter space above \( m_S = 1 \) TeV that is consistent with the XENON1T bound.

We remark the indirect signatures in the case of heavy scalar dark matter. For \( m_S > m_{LQ} \), dark matter can annihilate into a pair of leptoquarks, each of which decays into a pair of quark and lepton in cascade. In this case, the branching ratios of final products of DM annihilations are shown in Table 2. For \( m_S \gtrsim m_{LQ} \), a leptoquark pair is produced with almost zero velocities, so each leptoquark decays into a pair of quark and lepton such as \( \tilde{q}\tilde{q} \) or \( q'\tilde{l}' \), back-to-back. In this case, a pair of two quarks \( (q'q) \) or a pair of leptons \( (l'l) \) carry about the energy of DM mass, so we take them as if they are produced from the direct annihilations of dark matter with mass \( m_S/2 \) and impose the indirect detection bounds on the annihilation cross section. But, if \( m_S \gg m_{LQ} \), leptoquarks produced from the DM annihilations are boosted so the full energy spectra for quarks or leptons carry the energy spectra of wide box rather than a monochromatic energy. In this case, we need to take more care before imposing the indirect detection bounds.

When the boost effects of leptoquarks are ignored, for instance, \( m_{LQ} \gtrsim m_S \), we can apply the indirect detection bounds for the direct annihilations of dark matter to cascade annihilations. For instance, for a singlet leptoquark with \( \lambda_{33} \gg \lambda_{23} \) or \( \beta \approx 1 \), we get \( B(tt\tilde{\tau}) : B(\tilde{b}\tilde{b}\nu_{\tau}) : B(\tilde{t}\tilde{\tau}\nu_{\tau} + h.c.) = 1 : 1 : 1 \). Then, we can impose the Fermi-LAT diffuse gamma-ray constraint from \( b\bar{b} \) on \( \frac{1}{3}\langle \sigma v \rangle_{SS \to S_LQ S_LQ} \). Similarly, for a triplet leptoquark with \( \kappa_{32} \gg \kappa_{33} \) or \( \gamma \approx 1 \), we get \( B(\tilde{b}\tilde{b}\mu\nu_{\mu}) : B(\tilde{t}\tilde{t}\mu\nu_{\mu}) : B(\tilde{b}\tilde{b}\nu_{\mu}\nu_{\mu} + h.c.) : B(\tilde{t}\tilde{\tau}\nu_{\mu}\nu_{\mu}) = 1 : \frac{1}{3} : \frac{1}{3} : \frac{1}{3} : 1 \). In this case, we can impose the Fermi-LAT \( b\bar{b} \) bound on \( \frac{1}{5}\langle \sigma v \rangle_{SS \to S_LQ S_LQ} \) too. In general, positron, anti-proton and gamma-ray constraints are equally relevant for leptoquark-portal dark matter.
Figure 7: Relic density for scalar dark matter and various bounds in the parameter space, $\lambda_4$ vs $m_S$. The correct relic density can be obtained along the black and red lines, for models with singlet and triplet scalar leptoquarks, respectively. XENON1T bounds are shown in blue dashed lines. Indirect detection constraints from gamma-ray searches in Fermi-LAT (gray dotted) and HESS (brown dashed), and antiproton search in AMS-02 (pink dot-dashed) are overlaid. The bound from Higgs invisible decay is shown in purple dot-dashed line and the green regions are excluded by visible decays such as the Higgs diphoton signal strength.
Table 2: Branching ratios of products of DM annihilations into leptoquarks. Here, \( \beta \equiv \lambda_{33}^2 / (\lambda_{33}^2 + \lambda_{23}^2) \) and \( \gamma \equiv \kappa_{32}^2 / (\kappa_{32}^2 + \kappa_{22}^2) \).

As a consequence, the leptoquark-portal couplings lead to potentially distinct signatures with quarks and leptons mixed from the cascade annihilations of dark matter, as compared to the case with direct annihilations into a quark pair or a lepton pair. In other words, the region with \( m_S > m_{LQ} \) can be constrained by indirect detection experiments too. This interesting issue will be discussed in a future work.

5.4 Higgs data

The decay rate of the Higgs boson into a pair of dark matter particles is

\[
\Gamma(h \rightarrow SS) = \frac{\lambda_{33}^2 v^2}{32 \pi m_h} \sqrt{1 - \frac{4m_S^2}{m_h^2}}. \tag{56}
\]

The decay rate of the Higgs boson into a diphoton or a digluon is also modified due to leptoquarks, as given in eqs. (B.7) and (B.8). Corrections to \( h \rightarrow WW, ZZ \) are small because they are already present at tree level in the SM, so we can ignore them. Then, the total Higgs decay width is modified to

\[
\Gamma_h \approx \Gamma_{h, SM} + \Gamma(h \rightarrow SS) \tag{57}
\]

where \( \Gamma_{h, SM} = 4 \text{ MeV} \) in the SM. The bound from invisible Higgs decay, \( \text{BR}(h \rightarrow SS) < 0.24 \), leads to the following condition [58],

\[
\text{BR}(h \rightarrow SS) = \frac{\Gamma(h \rightarrow SS)}{\Gamma_h} < 0.24. \tag{58}
\]

The diphoton signal strength for gluon-fusion production is given by

\[
\mu_{\gamma\gamma} = R_{gg} R_{\gamma\gamma}, \tag{59}
\]
where

\[ R_{gg} = \frac{\sigma(gg \to h)}{\sigma(gg \to h)_{\text{SM}}} = \frac{\Gamma(h \to gg)}{\Gamma_h \cdot \text{BR}(h \to gg)_{\text{SM}}}, \quad R_{\gamma\gamma} = \frac{\Gamma(h \to \gamma\gamma)}{\Gamma_h \cdot \text{BR}(h \to \gamma\gamma)_{\text{SM}}}. \] (60)

The other visible decays, \( h \to ij \), are similarly modified to \( \mu_{ij} = R_{gg}R_{ij} \), through the modified total decay width of Higgs boson, with \( R_{ij} = \text{BR}(h \to ij)/\text{BR}(h \to ij)_{\text{SM}} = \Gamma_{h,\text{SM}}/\Gamma_h \). The measurements of \( gg \rightarrow h \rightarrow \gamma\gamma \) show \( \mu_{\gamma\gamma} = 1.10^{+0.23}_{-0.22} \) from the combined fit of LHC 7 TeV + 8 TeV data \[59\], and \( \mu_{\gamma\gamma} = 0.81^{+0.19}_{-0.18} \) and \( \mu_{\gamma\gamma} = 1.10^{+0.20}_{-0.18} \) from the ATLAS and CMS 13 TeV data, respectively \[60,61\].

In our model, as far as \( |\lambda_5| \lesssim 10 \), the decay rate into a diphoton or a digluon can be ignored, but the diphoton signal strength is modified by the enhanced total decay width of Higgs boson due to the invisible decay mode. This result can be read from Fig. 7 in the purple dot-dashed lines the region above which is excluded by Higgs invisible decay and in the green region which is excluded by the Higgs diphoton signal strength.

6 Conclusions

We have presented leptoquark models where scalar leptoquarks not only lead to the effective operators necessary for the \( B \)-meson anomalies but also become a portal to scalar dark matter through quartic couplings. We showed that the annihilations of dark matter into a leptoquark pair allow for a wide parameter space that is consistent with both the correct relic density and the XENON1T bound. These new annihilation channels lead to four-body final states in cascade with quarks and leptons mixed, due to the leptoquark decays. Therefore, there is an interesting interplay between the cascade annihilations of dark matter and the leptoquark search channels at the LHC, which can be tested in the current and future experiments.

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Appendix A: Effective Hamiltonians for \( B \)-meson decays.

From eq. (11), we obtain the relevant Yukawa couplings for the singlet scalar leptoquark
$S_1$ in components,

$$\mathcal{L}_{S_1} = -\lambda_{3j}\left((\bar{t}^C)_R S_1 l_{jL} - \langle b^C \rangle_R S_1 \nu_{jL}\right) + \text{h.c.}$$

$$-\lambda_{2j}\left((\bar{c}^C)_R S_1 l_{jL} - \langle s^C \rangle_R S_1 \nu_{jL}\right) + \text{h.c.} + \cdots. \quad (A.1)$$

Then, after integrating out the leptoquark $S_1$, we obtain the effective Hamiltonian relevant for $b \to c\tau\bar{\nu}_\tau$ as

$$\mathcal{H}_{b\to c\tau\bar{\nu}_\tau}^{S_1} = -\frac{\lambda_{33}^*\lambda_{23}}{m_{S_1}^2} (\bar{c}^C_R \gamma_\mu (b^C)_R) (\bar{\nu}_{\tau L} \gamma_\mu \tau_L) + \text{h.c.}$$

$$= -\frac{\lambda_{33}^*\lambda_{23}}{2m_{S_1}^2} \left(\bar{b}_L \gamma_\mu c_L\right) (\bar{\nu}_{\tau L} \gamma_\mu \tau_L) + \text{h.c.}. \quad (A.2)$$

where use is made of Fierz identity in the second line.

In particular, in MSSM, down-type squarks ($\tilde{b}_{Rk}^*$) \cite{62} belong to singlet scalar leptoquarks. We introduce the R-parity violating (RPV) superpotential as follows,

$$W \supset \lambda'_{ijk} L_i Q_j D_k^c,$$

resulting in the component field Lagrangian for doublet scalar leptoquarks $S_2 \equiv \tilde{u}_{Lk}$ with $Y = +\frac{1}{6}$ or singlet scalar leptoquarks $S_1 = \tilde{b}_{Rk}$ with $Y = +\frac{1}{3}$ as

$$\mathcal{L}_{\text{RPV}} = -\lambda'_{ijk} L_i \tilde{Q}_j d_k^c + \text{h.c.} + \cdots. \quad (A.4)$$

Picking up the necessary terms for $R_{K^0}$ and $R_{D^{(*)}}$ anomalies, we get, in terms of two component spinors,

$$\mathcal{L}_{\text{RPV}} = -\lambda'_{ijk} l_{jL} \tilde{u}_{Lk} b^c - \lambda'_{jk2} l_{jL} \tilde{u}_{Lk} s^c$$

$$-\lambda'_{j3k} \nu_{jL} b_l \tilde{b}_{Rk}^* - \lambda'_{j2k} l_{jL} c_L \tilde{b}_{Rk}^* + \text{h.c.} + \cdots. \quad (A.5)$$

Then, after integrating out the up-type squarks, $\tilde{u}_{Lk}$, and down-type squarks, $\tilde{b}_{Rk}^*$, we obtain the effective Hamiltonian for the semi-leptonic B-decays in terms of four-component spinors, as follows,

$$\mathcal{H}_{\text{eff}}^{\text{RPV}} = -\frac{\lambda'_{2k3}\lambda'_{2k2}}{m_{\tilde{u}_{Lk}}^2} (\bar{b}_{R\mu L})(\bar{\mu}_{Ls R})$$

$$-\frac{\lambda'_{32k}\lambda'_{33k}}{m_{d_{Rk}}^2} (\bar{c}_{R})(\bar{\mu}_{Ls R}) + \text{h.c.}$$

$$= -\frac{\lambda'_{2k3}\lambda'_{2k2}}{2m_{\tilde{u}_{Lk}}^2} (\bar{b}_{R\mu L})(\bar{\mu}_{Ls R}) - \frac{\lambda'_{32k}\lambda'_{33k}}{2m_{d_{Rk}}^2} (\bar{c}_{R})(\bar{\mu}_{Ls R}) + \text{h.c.}. \quad (A.6)$$

Therefore, the effective Hamiltonian for the $b$-to-$s$ transition is of the $(V + A)$ form, which was originally proposed to explain $R_K$ anomalies \cite{63,64} but is not consistent with $R_K$.  

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anomalies as it favors \((V - A)\) form. On the other hand, the effective Hamiltonian for the \(b\)-to-\(c\) transition is consistent with the \(R_{D(c)}\) anomalies \[30,65\].

From eq. (12), we obtain the relevant Yukawa couplings for the triplet leptoquark \(S_3\) in components,

\[
\mathcal{L}_{S_3} = -\kappa_{3j} \left( \sqrt{2} \left( t^C \right)_R \phi_3 \nu_{jL} - \left( t^C \right)_R \phi_2 \nu_{jL} - \sqrt{2} \left( b^C \right)_R \phi_2 l_{jL} + \right) + \text{h.c.}
\]

Then, after integrating out the leptoquark \(\phi_1\) with \(Q = \frac{4}{3}\), we obtain the effective Hamiltonian relevant for \(b \rightarrow s\mu^+\mu^-\) as

\[
\mathcal{H}_{b \rightarrow s\mu^+\mu^-}^{S_3} = -\frac{2\kappa^*_{32} \kappa_{22}}{m_{\phi_1}^2} \left( \left( s^C \right)_R \mu_L \left( b^C \right)_R \right) + \text{h.c.}
\]

\[
= -\frac{\kappa^*_{32} \kappa_{22}}{m_{\phi_1}^2} \left( \left( s^C \right)_R \gamma^\mu \left( b^C \right)_R \right) \left( \bar{\mu}_L \gamma^\mu \mu_L \right) + \text{h.c.}
\]

\[
= -\frac{\kappa^*_{32} \kappa_{22}}{m_{\phi_1}^2} \left( \bar{b}_L \gamma^\mu s_L \right) \left( \bar{\mu}_L \gamma^\mu \mu_L \right) + \text{h.c.}..\] (A.7)

Here, we note that use is made of the Fierz identity in the second line and \(\left( s^C \right)_R \gamma^\mu \left( b^C \right)_R = \bar{b}_L \gamma^\mu s_L\) is used in the third line.

The Yukawa couplings for the singlet scalar leptoquark also lead to effective Hamiltonian for \(b \rightarrow s\nu_i \bar{\nu}_j\) as follows,

\[
\mathcal{H}_{b \rightarrow s\nu_i \bar{\nu}_j}^{S_1} = \frac{\lambda^*_{3i} \lambda_{2j}}{m_{S_1}^2} \left( \left( s^C \right)_R \nu_{jL} \right) \left( \bar{\nu}_{iL} \left( b^C \right)_R \right) + \text{h.c.}
\]

\[
= \frac{\lambda^*_{3i} \lambda_{2j}}{2m_{S_1}^2} \left( \left( s^C \right)_R \gamma^\mu \left( b^C \right)_R \right) \left( \bar{\nu}_{iL} \gamma^\mu \nu_{jL} \right) + \text{h.c.}
\]

\[
= \frac{\lambda^*_{3i} \lambda_{2j}}{2m_{S_1}^2} \left( \bar{b}_L \gamma^\mu s_L \right) \left( \bar{\nu}_{iL} \gamma^\mu \nu_{jL} \right) + \text{h.c.}..\] (A.8)

A similar effective interactions can be obtained for the triplet scalar leptoquark, as discussed in the text.

**Appendix B: Effective interactions for dark matter and Higgs boson due to leptoquark loops.**

For heavy leptoquarks, we the effective interactions between scalar dark matter and SM gauge bosons, induced by leptoquarks, as follows,

\[
\mathcal{L}_{S,\text{eff}} = D_3 S^2 G_{\mu\nu} G^{\mu\nu} + D_2 S^2 W_{\mu\nu} W^{\mu\nu} + D_1 S^2 F_{\mu\nu} F^{\mu\nu} \] (B.1)
where

\[ D_3 = \frac{\alpha_S \lambda_3}{32 \pi m_{LQ}^2} N_{LQ} l_3(S_{LQ}) A_0(y), \quad (B.2) \]
\[ D_2 = \frac{\alpha \lambda_3}{32 \pi m_{LQ}^2} N_{c} l_2(S_{LQ}) A_0(y), \quad (B.3) \]
\[ D_1 = \frac{\alpha_Y \lambda_3}{32 \pi m_{LQ}^2} N_{c} N_{LQ} Y_{LQ}^2 A_0(y) \quad (B.4) \]

with

\[ A_0(y) = -y^{-2}[y - f(y)], \quad (B.5) \]

\[ f(y) = \begin{cases} \arcsin^2 \sqrt{y}, & y \leq 1, \\ -\frac{1}{4} \left[ \ln \frac{1+\sqrt{1-y^{-1}}}{1-\sqrt{1-y^{-1}}} - i\pi \right], & y > 1, \end{cases} \quad (B.6) \]

and \( y \equiv m_S^2 / m_{LQ}^2 \). Here, \( l_{2,3}(S_{LQ}) \) are the Dynkin indices of \( S_{LQ} \) under \( SU(2)_L \) and \( SU(3)_c \), respectively, i.e. \( l_3(S_{1,3}) = \frac{1}{2} \), \( l_2(S_1) = 0 \), and \( l_2(S_3) = 2 \), and \( N_{LQ} = 1, 3 \) for \( S_{LQ} = S_1, S_3 \), respectively.

Moreover, leptoquark couplings to the SM Higgs can modify the decay rates of Higgs boson into a diphoton or a digluon, as follows,

\[ \Gamma(h \to \gamma\gamma) = \frac{G_F \alpha_{em}^2 m_h^3}{128 \sqrt{2} \pi^3} \left| \sum_f N_c Q_f^2 A_{1/2}(x_f) + A_1(x_W) \right|^2 + N_c g_{LQ} \sum_{i=1,\ldots,N_{LQ}} Q_{LQ}^2 A_0(x_{LQ})^2, \quad (B.7) \]
\[ \Gamma(h \to gg) = \frac{G_F \alpha_{em}^2 m_h^3}{36 \sqrt{2} \pi^3} \left| \sum_f A_{1/2}(x_f) + \frac{3}{4} N_{LQ} g_{LQ} A_0(x_{LQ}) \right|^2 \quad (B.8) \]

where \( g_{LQ} \equiv \lambda_5 v^2 / (2m_{LQ}^2) \), \( x_i = m_i^2 / (4m_h^2) \) and the loop functions are

\[ A_{1/2}(x) = 2x^{-2}[x + (x - 1)f(x)], \quad (B.9) \]
\[ A_1(x) = -x^{-2}[2x^2 + 3x + 3(2x - 1)f(x)]. \quad (B.10) \]

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