Abstract

The observed muon anomalous magnetic moment deviates from the Standard Model (SM) predictions. There are two scalar leptoquarks with simultaneous couplings to the quark-muon pairs of both chiralities that can singly explain this discrepancy. We discuss an alternative mechanism that calls for mixing of two scalar leptoquarks of the same electric charge through the interaction with the Higgs field, where the two leptoquarks separately couple to the quark-muon pairs of opposite chirality structures. Three scenarios that satisfy this requirement are $S_1 \& S_3$, $\tilde{S}_1 \& S_3$, and $\tilde{R}_2 \& R_2$, where the first scenario is realised with the up-type quarks running in the loops while the other two scenarios proceed through the down-type quark loops. We constrain the leptoquark mixing parameters with oblique corrections and introduce only two non-zero Yukawa couplings to quarks and muon, at the time, to study ability of these three scenarios to explain $(g - 2)_\mu$ and be in accord with existing constraints. We find that the $S_1 \& S_3$ scenario with the (charm) top quark loops is (not) viable, whereas the $\tilde{S}_1 \& S_3$ and $\tilde{R}_2 \& R_2$ scenarios require at least one of the two Yukawa couplings to be an $O(1)$ parameter to accommodate the $(g - 2)_\mu$ discrepancy. If Yukawa couplings are to remain perturbative for the $S_1 \& S_3$ scenario with the top quark loops, we find an upper bound on the leptoquark mass scale to be at 15 TeV.
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1 Introduction

Leptoquarks (LQs) are hypothetical particles of either scalar or vector nature that couple at tree-level to quark-lepton pairs [1, 2]. The fact that they carry non-trivial baryon (B) and lepton (L) numbers makes them particularly appealing sources of New Physics (NP) regarding phenomena related to the B and/or L number violation. These phenomena include proton decay [3] and generation of neutrino mass [4, 5] to name a few.

In this work we are interested in the B and L number conserving effects of scalar LQs with regard to the anomalous magnetic moment of muon. The main reason behind this study is the long-standing discrepancy between the measured value [6] and theoretical predictions of that observable [7–9]. The experimental result \( a_{\mu}^{\exp} \) for the muon anomalous magnetic moment deviates from the Standard Model (SM) predictions \( a_{\mu}^{\text{SM}} \) roughly at the level of 4 \( \sigma \) [10]. More precisely, the discrepancy that we want to address currently reads

\[
\delta a_{\mu} = a_{\mu}^{\exp} - a_{\mu}^{\text{SM}} = (2.7 \pm 0.8) \times 10^{-9}.
\]

Our interest is further amplified by the anticipated experimental result from the Muon g–2 Collaboration [11] at Fermilab that might clarify the nature of the current disagreement.

The influence of scalar LQs on \( a_{\mu} = (g - 2)_{\mu}/2 \) is well-documented in the literature [12]. The only contribution that might be large enough to address the observed difference is of the one-loop nature and it requires a presence of a non-chiral LQ [13]. The non-chiral LQs are those that couple to both left- and right-chiral quarks of the same type and, due to Lorentz symmetry, to the charged leptons of opposite chiralities, where we assume that the fermion content of the NP scenario is purely the SM one. It turns out that the only scalar LQ multiplets that possess non-chiral couplings to muons are \( R_2 \) and \( S_1 \). We specify, for completeness, transformation properties of all scalar LQs under the SM gauge group \( SU(3)_c \times SU(2)_L \times U(1)_Y \) and associated nomenclature in Table 1. We also show the generic one-loop Feynman diagrams induced by the non-chiral scalar LQs that yield contributions to \((g - 2)_{\mu}\) in Fig. 1. Clearly, the SM fermions that close the \((g - 2)_{\mu}\) loop, in the non-chiral LQ case, are the up-type quarks while the leading NP contribution is (inversely) proportional to the mass of the (LQ) quark in the loop.

| Symbol | \((SU(3)_c, SU(2)_L, U(1)_Y)\) | Interactions | \( F = 3B + L \) |
|--------|---------------------------------|--------------|-----------------|
| \( S_3 \) | \((3, 3, 1/3)\) | \( \overline{Q}^C L \) | -2 |
| \( R_2 \) | \((3, 2, 7/6)\) | \( \overline{u}_R L, \overline{Q}_e R \) | 0 |
| \( \tilde{R}_2 \) | \((3, 2, 1/6)\) | \( \overline{d}_R L \) | 0 |
| \( \tilde{S}_1 \) | \((3, 1, 4/3)\) | \( \overline{d}_R^C e_R \) | -2 |
| \( S_1 \) | \((3, 1, 1/3)\) | \( \overline{Q}^C L, \overline{\pi}_R^C e_R \) | -2 |

Table 1: List of all scalar LQs, their SM quantum numbers, renormalizable interactions to the quark-lepton pairs, and associated fermion numbers. Interactions with right-handed neutrinos are not considered.
In this work we investigate the viability of those scenarios where the one-loop contributions towards the anomalous magnetic moment of muon are induced through the mixing of two scalar LQs of the same electric charge via the SM Higgs field, where the LQs in question need to couple to the muons of opposite chiralities. We accordingly study the existing constraints on the parameter space of this particular mechanism due to electroweak precision measurements, relevant flavor observables, and the current LHC analyses.

The paper is organized as follows. In Sec. 2, we describe single LQ contributions to $(g - 2)_{\mu}$ to set the notation. In Sec. 3, we classify those pairs of scalar LQs that can mix via the SM Higgs field and subsequently generate chirality-enhanced contributions towards $(g - 2)_{\mu}$. We find three possible LQ pairs – $S_1 \& S_3$, $\tilde{S}_1 \& S_3$, and $R_2 \& R_2$ – that might generate large enough contributions towards $(g - 2)_{\mu}$ through the mixing with the SM Higgs field. We proceed to discuss electroweak precision constraints on the LQ mixing and discuss relevant differences between the three scenarios, if any. We then confront, in Sec. 4, the $S_1 \& S_3$ scenario with $(g - 2)_{\mu}$ and various phenomenological constraints to investigate its viability. The ability of the $\tilde{S}_1 \& S_3$ and $\tilde{R}_2 \& R_2$ scenarios to address $(g - 2)_{\mu}$ is briefly discussed in Secs. 5 and 6, respectively. We summarize our findings in Sec. 7.

2 Single LQ contributions to $(g - 2)_{\mu}$

The most general formulae for the interactions of the generic scalar LQ $S$ of the definite fermion number $F = 3B + L$ with the quark-charged lepton $(q-\ell)$ pairs, in the mass eigenstate basis, are [12]

\[
\mathcal{L}^F=0 = \overline{q}_i (l^{ij} P_R + r^{ij} P_L) \ell_j S + \text{h.c.} \\
\mathcal{L}^{|F|=2} = \overline{q}_i^C (l^{ij} P_L + r^{ij} P_R) \ell_j S + \text{h.c.},
\]

(2)

where $i$ ($j$) represents generation index for quarks (charged leptons), $P_{L,R}$ are projection operators, and $l^{ij}$ and $r^{ij}$ are Yukawa coupling strengths. The LQ contributions to the muon anomalous moment can then be written in the compact form [2]

\[
\delta a_{\mu} = -\frac{N_c m_{\mu}}{8\pi^2 m_S^2} \sum_q \left[ m_{\mu} (|l^{\mu q}|^2 + |r^{\mu q}|^2) F_{QS}(x_q) + m_q \Re (r^{\mu q} l^{\mu q}) G_{QS}(x_q) \right],
\]

(3)
where \( m_\mu \) is muon mass, \( m_S \) \( (m_q) \) is the LQ (quark) mass, \( x_q = m_q^2/m_S^2 \), \( N_c = 3 \) is the number of colors, and

\[
F_{Qs}(x) = Qs f_S(x) - f_F(x),
\]

\[
G_{Qs}(x) = Qs g_S(x) - g_F(x).
\]

We denote with \( Qs \) the electric charge of scalar \( S \) and define the loop functions

\[
f_S(x) = \frac{x + 1}{4(1 - x)^2} + \frac{x \log x}{2(1 - x)^3}, \quad g_S(x) = \frac{1}{x - 1} - \frac{\log x}{(x - 1)^2},
\]

\[
f_F(x) = \frac{x^2 - 5x - 2}{12(x - 1)^3} + \frac{x \log x}{2(x - 1)^4}, \quad g_F(x) = \frac{x - 3}{2(x - 1)^2} + \frac{\log x}{(x - 1)^3}.
\]

It is clear from Eq. (3) that the scalar LQs that have only left- or right-chiral couplings generate contributions to \( a_\mu \) that are of a definite sign and that are suppressed by the lepton mass. One thus minimally needs either one non-chiral LQ or two scalar LQs of the same electric charge that can mix through the Higgs field and that can couple to muons of opposite chiralities to generate phenomenologically viable shift in \( a_\mu \). Again, according to the content of Table 1, the only LQs that can potentially generate Feynman diagrams shown in Fig. 1 are \( R_2 \) and \( S_1 \) with \( q \in \{u,c,t\} \).

There exists a number of dedicated studies of a single LQ resolution of the muon anomalous magnetic moment discrepancy [14–18]. We, in this manuscript, turn our attention to the explanation of \( (g - 2)_\mu \) within the LQ mixing scenario context that we define next.

## 3 LQ mixing scenarios

### 3.1 General classification

We list in Table 2 all possible pairs of scalar LQs that can mix through the SM Higgs field [19]. We require that the mixing term is exclusively due to an interaction between the two LQs and one or, at most, two SM Higgs fields. We denote the SM Higgs field with \( H = (1,2,1/2) \) and indicate the number of \( H \) fields in the contraction in the second column of Table 2. Those LQ pairs that couple to muons of opposite chiralities can generate the one-loop contribution towards \( (g - 2)_\mu \) that we are interested in whereas the pairs that couple to neutrinos can yield neutrino masses of Majorana nature [20–22] at the one-loop level. In both instances the LQ pairs need to couple to the quarks of the same type but opposite chiralities in order for the relevant loop(s) to be completed and, in case of \( (g - 2)_\mu \), for associated contributions to be enhanced. We accordingly specify the quark type that is in the loop, where we use \( u \) and \( d \) to collectively denote the up-type and the down-type quarks, respectively.
Table 2: Scalar LQ pairs that can, through the mixing with the SM Higgs field, generate either the one-loop contributions towards $(g-2)_\mu$ or neutrino mass. It is indicated whether the chirality-enhanced contributions are proportional to the up-type ($u$) or down-type ($d$) quark masses.

| LQ pairs         | Mixing field(s) | $(g-2)\mu$ | $\nu$-mass |
|------------------|-----------------|-------------|-------------|
| $S_1 & S_3$      | $H \ H$        | $u$         | $-$         |
| $\tilde{S}_1 & S_3$ | $H \ H$        | $d$         | $-$         |
| $\tilde{R}_2 & R_2$ | $H \ H$        | $d$         | $-$         |
| $\tilde{R}_2 & S_1$ | $H$             | $-$         | $d$         |
| $\tilde{R}_2 & S_3$ | $H$             | $-$         | $d$         |

There are, clearly, three possible LQ pairs that might generate large enough contributions towards $(g-2)_\mu$ through the mixing with the SM Higgs field. These combinations are $S_1 & S_3$, $\tilde{S}_1 & S_3$, and $\tilde{R}_2 & R_2$, where, in all three instances, at least one of the LQ multiplets is chiral in nature. More importantly, at least one of the two LQ multiplets that mix carries non-trivial $SU(2)_L$ assignment. The LQ mixing mechanism can thus induce mass splitting between the states belonging to the same LQ multiplet that, in turn, might generate substantial oblique corrections that could be in conflict with the existing electroweak precision measurements.

### 3.2 Mixing formalism

To describe the most prominent features of the LQ mixing we assume existence of two scalars $S_a^{(Q)}$ and $S_b^{(Q)}$, of the same electric charge $Q$, but from two different multiplets $S_a$ and $S_b$ that might have non-trivial weak isospins $I^{S_a}$ and $I^{S_b}$. We thus expect, on general grounds, to have $2(I^{S_a} + I^{S_b} + 1)$ mass eigenstates with or without mixing. The mass squared matrix for the mixed states reads

$$\mathcal{M}^2 = \begin{pmatrix} m_{S_a}^2 & \Omega \\ \Omega & m_{S_b}^2 \end{pmatrix},$$

(8)

where $m_{S_a}$ and $m_{S_b}$ denote the common masses of all $S_a$ and $S_b$ components prior to the mixing and $\Omega$ stands for the mixing term arising from the interactions of $S_a$ and $S_b$ with the Higgs boson that we discuss later in concrete scenarios. The matrix in Eq. (8) can be brought into diagonal form with a simple field redefinition

$$\begin{pmatrix} S_{-}^{(Q)} \\ S_{+}^{(Q)} \end{pmatrix} = \begin{pmatrix} \cos \theta & \sin \theta \\ -\sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} S_a^{(Q)} \\ S_b^{(Q)} \end{pmatrix},$$

(9)

where $S_{\pm}^{(Q)}$ are the mass eigenstates and the mixing angle $\theta \in [-\pi/2, \pi/2]$ is related to $\Omega$ via the relation

$$\tan 2\theta = \frac{2 \Omega}{m_{S_a}^2 - m_{S_b}^2}.$$  

(10)
The physical masses squared for the mixed states are

\[
m^2_{S_{\pm Q}} = \frac{m_{S_a}^2 + m_{S_b}^2}{2} \pm \frac{1}{2} \sqrt{(m_{S_a}^2 - m_{S_b}^2)^2 + 4\Omega^2}.
\]

(11)

where for simplicity we assume that \(m_{S_b} \geq m_{S_a}\). There are additional \(2I^{S_a} (2I^{S_b})\) mass eigenstates in multiplet \(S_a (S_b)\) with masses \(m_{S^a} (m_{S^b})\) besides the two mass eigenstates \(S_{\pm}^{(Q)}\).

The LQs mix maximally, according to Eq. (10), for \(m_{S_a} = m_{S_b} \equiv m_S\) with \(\theta = \pi/4\). In this particular case, if we take the limit when \(|\Omega| \ll m_{\tilde{S}}^2\), it is convenient to define the parameter

\[
\delta m_{S}^{(Q)} = m_{S^{(Q)}} - m_S \approx m_S - m_{S^{(Q)}} ,
\]

(12)

which is directly related to the strength of the LQ mixing. Note that \(|\Omega|/m_S^2\) is indeed a small parameter due to the actual form of \(\Omega\) in our set-up and the current limits from LHC on LQ masses as we explicitly show below.

The preceding expressions can now be straightforwardly applied to the three LQ combinations listed in Table 1.

- **\(S_1 \& S_3\):** The interactions of \(S_3 = S_a\) and \(S_1 = S_b\) with the Higgs boson \(H\) read

  \[
  \mathcal{L}_{\text{mix}}^{S_1 \& S_3} = \xi H^{\dagger}(\vec{\tau} \cdot \tilde{S}_3)HS_1^* + \text{h.c.,}
  \]

  (13)

  where \(\xi\) is a dimensionless coupling that, after electroweak symmetry breaking, induces a mixing between the \(Q = 1/3\) states \(S_3^{(1/3)}\) and \(S_1^{(1/3)}\). The mass eigenstates of the \(Q = 1/3\) fields are then described by Eq. (9) for \(S_a^{(Q)} = S_3^{(1/3)}\), \(S_b^{(Q)} = S_1^{(1/3)}\), and \(\Omega = -\xi v^2/2\). Since the weak isospins of LQ multiplets \(S_1\) and \(S_3\) are \(I^{S_1} = 0\) and \(I^{S_3} = 1\), respectively, this scenario has four mass eigenstates that are \(m_{S_{\pm}^{(1/3)}}\) and \(m_{S_3} \equiv m_{S_{3}^{(4/3)}} = m_{S_{3}^{(-2/3)}}\).

- **\(\tilde{S}_1 \& \tilde{S}_3\):** The interactions of \(S_3 = S_a\) and \(\tilde{S}_1 = S_b\) with \(H\) are described by

  \[
  \mathcal{L}_{\text{mix}}^{\tilde{S}_1 \& \tilde{S}_3} = \xi H^{\dagger}i\tau_2(\vec{\tau} \cdot \tilde{S}_3)HS_1^* + \text{h.c.,}
  \]

  (14)

  which induce a mixing between the \(Q = 4/3\) component of \(S_3\) with \(\tilde{S}_1\), where \(S_a^{(Q)} = S_3^{(4/3)}\) and \(S_b^{(Q)} = \tilde{S}_1^{(1/3)}\). The parameter \(\xi\) is normalized in such a way that, again, one gets \(\Omega = -\xi v^2/2\) in Eq. (8). The physical masses in this particular scenario are \(m_{S_{\pm}^{(4/3)}}\) and \(m_{S_3} = m_{S_{3}^{(1/3)}} = m_{S_{3}^{(-2/3)}}\).

- **\(\tilde{R}_2 \& R_2\):** The \(\tilde{R}_2 \& R_2\) pair mixes through the operator (see also Ref. [23])

  \[
  \mathcal{L}_{\text{mix}}^{\tilde{R}_2 \& R_2} = -\xi \left(R_2^TH\right)\left(\tilde{R}_2^T\tau_2H\right) + \text{h.c.}
  \]

  (15)
mixing. We analyse all three LQ mixing scenarios in what follows.

We shall now discuss the constraints on the mixing parameter \( \xi \), or equivalently on the mass splitting parameter \( \delta m_S^{(Q)} \) of Eq. (12) for all three scenarios, i.e., when \( Q = 1/3, 4/3, 2/3 \).

### 3.3 Electroweak constraints

The most relevant constraint on the LQ mixing arises from the electroweak precision data and, most particularly, the oblique parameter \( T \), which is known to constrain the mass splitting within a given electroweak multiplet. As we will show shortly, this constraint is crucial in estimating the largest allowed contribution towards \((g - 2)_\mu\) coming from LQ mixing. We analyse all three LQ mixing scenarios in what follows.

- **\( S_1 \& S_3 \):** First, we consider the \( S_1 \& S_3 \) scenario, which is described by the masses \( m_{S_3} \equiv m_{S_3(-2/3)} = m_{S_3(4/3)} \) and \( m_{S_3(1/3)} \). We find that the modification of the \( T \)-parameter we define to be \( \Delta T = T - T^{\text{SM}} \), in this particular scenario, is

  \[
  \Delta T_{S_1 \& S_3} = \frac{N_c}{4\pi s_W^2 m_W^2} \left[ \cos^2 \theta F(m_{S_3}, m_{S_{3(-2/3)}}) + \sin^2 \theta F(m_{S_3}, m_{S_{3(1/3)}}) \right],
  \]

  with

  \[
  F(m_1, m_2) = m_1^2 + m_2^2 - \frac{2m_1^2 m_2^2}{m_1^2 - m_2^2} \log \left( \frac{m_1^2}{m_2^2} \right),
  \]

where we introduce the shorthand notation \( \sin \theta_W = s_W \), with \( \theta_W \) being the Weinberg angle. In the degenerate case, i.e., when \( m_{S_1} = m_{S_3} \equiv m_S \), the expression in Eq. (16) can be expanded in terms of a mass splitting parameter \( \delta m_S^{(1/3)} \equiv m_S - m_{S_{3(-2/3)}} \approx m_{S_{3(1/3)}} - m_S > 0 \), giving at leading order,

\[
\Delta T_{S_1 \& S_3} = \frac{N_c}{3\pi s_W^2} \left( \frac{\delta m_S^{(1/3)}}{m_W} \right)^2 \left[ 1 + \mathcal{O} \left( \frac{\delta m_S^{(1/3)}}{m_S} \right)^2 \right],
\]

By using the inputs from the most recent electroweak fit, i.e., \( \Delta T = 0.05(12) \) \cite{24}, we obtain the following 2 \( \sigma \) constraint,

\[
\left[ \delta m_S^{(1/3)} \right]_{S_1 \& S_3} \lesssim 40 \text{ GeV}.
\]

- **\( \tilde{S}_1 \& S_3 \):** Similar conclusions apply to the \( \tilde{S}_1 \& S_3 \) scenario, in which we find that \( \Delta T \) is given by Eq. (16) that is divided by a factor of 2, with the appropriate replacement of the mixing parameters, as discussed in Sec. 3.2. The factor of 2 comes from the
fact that the $I_3^{S_3} = +1$ component of $S_3$ mixes with $S_1$, while in the $S_1 \& S_3$ case the mixing is between the $I_3^{S_3} = 0$ component of $S_3$ and $S_1$. The LEP constraint, in the $\tilde{S}_1 \& S_3$ case, is thus slightly weaker. Namely, in the maximal-mixing scenario, we find

$$\left[ \delta m_S^{(4/3)} \right]_{S_1 \& S_3} \lesssim 50 \text{ GeV}. \quad (20)$$

- $\tilde{R}_2 \& R_2$: Lastly, we discuss electroweak constraints for the $\tilde{R}_2 \& R_2$ scenario. In this case, we find that the operator introduced in Eq. (15) induces the following contribution to $\Delta T$,

$$\Delta T_{\tilde{R}_2 \& R_2} = \frac{N_c}{16\pi s_W^2} \frac{1}{2} \left\{ \cos^2 \theta \left[ F(m_{R_2}, m_{S_2^{(2/3)}}) + F(m_{\tilde{R}_2}, m_{S_2^{(2/3)}}) \right] \right.$$

$$+ \sin^2 \theta \left[ F(m_{R_2}, m_{S_2^{(2/3)}}) + F(m_{\tilde{R}_2}, m_{S_2^{(2/3)}}) \right] \left. \right\}, \quad (21)$$

where $m_{R_2} \equiv m_{R_2^{(5/3)}}$ and $m_{\tilde{R}_2} \equiv m_{\tilde{R}_2^{(-1/3)}}$, while $m_{S_2^{(2/3)}}$ are described by Eq. (11) with the appropriate replacements. In the maximal mixing case, i.e., for $m_{R_2} = m_{\tilde{R}_2} \equiv m_S$, we define $\delta m_{S}^{(2/3)} \equiv m_S - m_{S_2^{(2/3)}} \approx m_{S_2^{(2/3)}} - m_S$, so that

$$\Delta T_{\tilde{R}_2 \& R_2} = \frac{N_c}{6\pi s_W^2} \left( \frac{\delta m_{S}^{(2/3)}}{m_W} \right)^2 \left[ 1 + O \left( \frac{\delta m_{S}^{(2/3)}}{m_S} \right)^2 \right]. \quad (22)$$

By reinterpreting LEP data, we then find that the mass splitting is constrained to

$$\left[ \delta m_{S}^{(2/3)} \right]_{\tilde{R}_2 \& R_2} \lesssim 50 \text{ GeV}, \quad (23)$$

similarly to the $\tilde{S}_1 \& S_3$ scenario.

With these ingredients at hands, we will now discuss the contributions to $(g - 2)_\mu$ in each of the potentially viable scenarios.

4 \hspace{1em} $(g - 2)_\mu$ via $S_1 \& S_3$

4.1 Setup

The $S_1 \& S_3$ scenario is the only scenario for which top quark runs in the $(g - 2)_\mu$ loops and it is thus the most promising one of the three. The relevant interactions of $S_1$ and $S_3$ with the SM fermions are given by

$$\mathcal{L}_{S_1} = y^{ij}_{R} \bar{u}_{R_i} \tau_{S_j} S_1 + h.c., \quad (24)$$

$$\mathcal{L}_{S_3} = y^{ij}_{L} \bar{Q}_i \tau S_3 L_j + h.c., \quad (25)$$

where $y_R$ and $y_L$ are Yukawa coupling matrices and $i$ and $j$ are flavor indices for quarks and leptons, respectively. We omit $B$ number violating couplings of both $S_1$ and $S_3$ as
well as the couplings of $S_1$ with the left-chiral leptons. The latter are omitted since we want to investigate the chirality-enhanced contributions coming solely from the LQ mixing scenario. The Lagrangians in Eqs. (24) and (25) can be expanded in terms of the electric charge eigenstates as

$$
\mathcal{L}_{S_1, S_3} = y_{ij}^R \bar{\nu}^C_R e_R j \bar{S}_1^{(1/3)} - y_{ij}^L \bar{\nu}^C_L j L \bar{S}_3^{(1/3)} - \sqrt{2} y_{ij}^L \bar{\nu}^C_L j L \bar{S}_3^{(4/3)} + \sqrt{2} (V^* y_{ij}^L) \bar{\nu}^C_L j L \bar{S}_3^{(-2/3)} - (V^* y_{ij}^L) \bar{\nu}^C_L j L \bar{S}_3^{(1/3)} + \text{h.c. ,}
$$

(26)

where $V$ is a Cabibbo-Kobayashi-Maskawa mixing matrix. We take a Pontecorvo-Maki-Nakagawa-Sakata mixing matrix to reside in neutrino sector. In this convention $\nu_{Lj}$ are flavor eigenstates. Note that, after electroweak symmetry breaking, one should replace $S_3^{(1/3)}$ and $S_1^{(1/3)}$ in Eq. (26) with $S_{\pm} S_{\mp}$, as defined in Eqs. (9) and (13), to go to the mass eigenstate basis for electrically charged fields.

LQ mixing can induce chirality-enhanced contributions to $(g - 2)_\mu$, as illustrated in Fig. 2. To compute $\delta a_\mu$ within this scenario, we replace each of the mass eigenstates running in the loops, i.e., $S_3^{(4/3)}$ and $S_{\pm}$, and their Yukawa interactions in Eq. (3). Note that we assume that $y_{ij}^R$ and $y_{ij}^L$, as defined in Eqs. (24) and (25), are the only non-zero entries in $y_R$ and $y_L$. The final result reads

$$
\delta a_\mu = -\frac{N_c m_{\mu}^2}{8\pi^2} \left\{ 2 |y_L| |y_R| \frac{\mathcal{F}_{4/3}(x_b)}{m_{S_3}^2} + \left[ \sin^2 \theta |y_{L_b}^\mu|^2 + \cos^2 \theta |y_{L^\mu}^b|^2 \right] \frac{\mathcal{F}_{4/3}(x_b)}{m_{S_3}^2} + \left[ \cos^2 \theta |y_{L_b}^\mu|^2 + \sin^2 \theta |y_{L^\mu}^b|^2 \right] \frac{\mathcal{F}_{4/3}(x_t)}{m_{S_+}^2} + \frac{m_t}{m_\mu} \sin \theta \cos \theta \Re (y_{L_b}^\mu y_{L^\mu}^b) \left[ \frac{\mathcal{G}_{4/3}(x_b)}{m_{S_+}^2} - \frac{\mathcal{G}_{4/3}(x_t)}{m_{S_-}^2} \right] \right\},
$$

(27)

where $x_b = m_{b_b}^2/m_{S_3}$ and $x_t^\pm = m_t^2/m_{S_{\pm}}$. These expressions have been further simplified by taking $V_{tb} \approx 1$. The loop functions $\mathcal{F}_Q$ and $\mathcal{G}_Q$ are defined in Eqs. (4) and (5). In the absence of LQ mixing, the chirality-enhanced contribution in the last line would vanish since $m_{S_+} = m_{S_-}$.

One particularly simple realization of the $S_1 \& S_3$ scenario concerns the case when the LQs mix maximally, i.e., for $\theta \approx \pi/4$, that arises when $m_{S_3} \approx m_{S_1}$. In this case we define $m_S$ and $\delta m_S^{(1/3)}$ as in Eq. (12), and expand Eq. (27) as a series in $\delta m_S^{(1/3)}/m_S$. We obtain,
to leading order, that

\[
\left[ \delta a_\mu \right]_{\theta = \pi/4} \simeq - \frac{N_c m_\mu^2}{4 \pi^2 m_S^2} \left\{ 2 |y_L^{\mu}|^2 \mathcal{F}_{4/3}(x_b) + (|y_L^{b\mu}|^2 + |y_R^{b\mu}|^2) \mathcal{F}_{1/3}(x_t) - 4 \frac{\delta m_S^{(1/3)}}{m_S} \frac{m_t}{m_\mu} \text{Re}(y_L^{b\mu} y_R^{t\mu}) [\mathcal{G}_{1/3}(x_t) + x_t \mathcal{G}'_{1/3}(x_t)] \right\} ,
\]

where \( x_t = m_t^2/m_S^2 \) and

\[
\mathcal{G}_{1/3}(x_t) + x_t \mathcal{G}'_{1/3}(x_t) = - \frac{7 x_t + 11}{6(x_t - 1)^3} + \frac{x_t^2 + 6 x_t + 2}{3(x_t - 1)^4} \log x_t .
\]

We recognize in the first line of Eq. (28) the chirality conserving contribution while the chirality-flipping one appears in the second line. The latter is proportional to \( \delta m_S^{(1/3)}/m_S \) and, as expected, vanishes in the limit of negligible mixing, i.e., when \( \delta m_S^{(1/3)} \approx 0 \). Clearly, electroweak precision tests represent a useful handle to constrain \( \delta m_S^{(1/3)} \) and to estimate the dominant contribution to \( (g - 2)_\mu \) in this scenario.

### 4.2 Phenomenological constraints

We shall now discuss the constraints on the LQ Yukawa couplings. Again, to obtain the largest possible contribution to \( \delta a_\mu \), coming from top quark in the loops, we assume that \( y_R^{t\mu} \) and \( y_L^{b\mu} \) of Eqs. (24) and (25) are the only non-zero Yukawa couplings. With this ansatz, most of the low-energy observables are unaffected at first approximation by the LQ presence. The only exceptions are the \( Z \)-pole observables and the LHC signatures, as we describe below.

- **\( Z \rightarrow \ell\ell \)**: LQs contribute to the \( Z \)-boson couplings to leptons at the one-loop level. These contributions can be generically parameterized as

\[
\mathcal{L}_{\text{eff}}^Z = \frac{g}{c_W} \sum_{f, i, j} \bar{f}_i \gamma^\mu \left[ g_{ij}^{fL} P_L + g_{ij}^{fR} P_R \right] f_j Z^\mu ,
\]

where \( g \) is the \( SU(2)_L \) gauge coupling and \( g_{ij}^{fL(R)} = \delta_{ij} g_{f(L(R)}^{SM} + \delta g_{f(L(R)}^{ij} \) with \( g_{fL}^{SM} = I_3^f - Q^f s_W^2 \) and \( g_{fR}^{SM} = -Q^f s_W^2 \), while \( \delta g_{f(L(R)}^{ij} \) contains the LQ contributions. In these expressions, \( I_3^f \) denotes the third weak-isospin component and \( Q_f \) denotes the electric charge of a fermion \( f \in \{ u, d, \ell, \nu \} \). The generic LQ contributions to \( \delta g_{f(L(R)}^{ij} \), at the one-loop order, have been computed in Ref. [25] and are incorporated in our analysis. For the experimental inputs, we consider the fit to LEP data performed in Ref. [26], which includes the correlation between the \( W \) and \( Z \) boson couplings to leptons and neutrinos coming from gauge invariance.

- **The LHC searches**: LQ \( S \) can be pair-produced in hadron colliders via \( gg(q\bar{q}) \rightarrow S^* S \). These pairs can be searched for via subsequent decay into quark-lepton pairs.

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See, for example, Refs. [27, 28] for recent reviews on the subject. For the scenarios we consider, the most stringent bounds come from direct searches for scalar LQs decaying into $b\mu$ ($t\mu$) final states, setting the limit $m_S \gtrsim 1400$ GeV ($m_S \gtrsim 1420$ GeV) for $\mathcal{B}(S \to b\mu) = 1$ [27] ($\mathcal{B}(S \to t\mu) = 1$ [28]). If we considered instead an LQ coupling to second generation quarks, the limit obtained would be $m_S \gtrsim 1530$ GeV for $\mathcal{B}(S \to j\mu) = 1$ [29]. In our phenomenological analysis we mainly investigate the scenarios that introduce four, for all practical purposes, mass degenerate LQ states, some of which certainly have the same decay signatures to final states with muons. To accommodate this, we use the approach advocated in Ref. [27] to deduce that the conservative mass limit, using experimental input of Refs. [30, 31], is $m_S \gtrsim 1.6$ TeV and we use it in what follows.

LQ couplings can also be indirectly probed via the study of high-$p_T$ dilepton-tails at LHC [32]. In this way, one can constrain, for instance, the LQ coupling $y_{L\mu}^{b\mu}$ as a function of LQ mass. In our analysis, we consider the constraints obtained in Ref. [28] for different LQ states with couplings to muons from 36 fb$^{-1}$ of LHC data and offer projection for the 300 fb$^{-1}$ data set.

4.3 Numerical results

Maximal mixing

We start by considering the $S_1$ & $S_3$ scenario with maximal mixing, i.e., when $\theta = \pi/4$, and with top quark in the $(g-2)\mu$ loops. In this particular case, $m_{S_1} = m_{S_3} \equiv m_S$ and we are left with four parameters

$$\{m_S, \delta m_S^{(1/3)}, y_{L\mu}^{b\mu}, y_{R\mu}^{t\mu}\}$$

(31)

that are defined in Eqs. (12), (24), and (25). We choose the mass splitting parameter $\delta m_S^{(1/3)}$ to saturate the bound obtained in Eq. (19) from oblique parameters and we impose the constraints from Sec. 4.2 onto the LQ Yukawa couplings. The parameter space of $y_{L\mu}^{b\mu}$ vs. $y_{R\mu}^{t\mu}$ that can address $(g-2)\mu$ is then shown in Fig. 3 for a benchmark mass $m_S = 1.6$ TeV. It is clear from Fig. 3 that the Yukawa couplings that are needed to explain $(g-2)\mu$ can be rather small and in full agreement with all the existing constraints.

Next, we allow $m_S$ to vary and we estimate its maximal value that is still consistent with the perturbative bound on Yukawa couplings, where we conservatively take $|y_{L\mu}^{b\mu}|, |y_{R\mu}^{t\mu}| \lesssim \sqrt{4\pi}$. We find that the maximal allowed mass is given by $m_S \lesssim 15$ TeV that is possibly within reach of high-energy LHC. Note that this limit is considerably lower than the one obtained for the single LQ solutions of the $(g-2)\mu$ discussed in Sec. 2. This can be understood from the additional suppression in Eq. (28), by a factor $\delta m_S^{(1/3)}/m_S \approx 10^{-2}$, that comes from the mixing itself.

General case

A natural question is whether the conclusions drawn above are changed in a scenario where the LQs do not mix maximally, as we shall consider now. In the general case, the
Figure 3: The values of $y^\mu_{tR}$ vs. $y^\mu_{bL}$ for the $S_1 \& S_3$ scenario that satisfy relevant flavor constraints and address $(g-2)_\mu$ through the top quark loops are shown in dark (light) blue to 1 σ (2 σ) accuracy for $m_{S_1} = m_{S_3} = 1.6$ TeV and $\delta m_S^{1/3} = 40$ GeV. The individual 2σ constraints from $Z$-pole observables and $(g-2)_\mu$ are represented with green and red dashed lines, respectively. The (projected) LHC limits from the study of $pp \to \mu\mu$ at high-$p_T$, at the 2σ level, for (300 fb$^{-1}$) 36 fb$^{-1}$ worth of data are shown with (dash-dotted) dashed gray lines [28].

The mixing angle $\theta$ is then determined by Eq. (10), as a function of $m_{S_{\pm}}$ and $\xi$. To answer this question, we fix the masses to benchmark values $m_{S_-} = 1.6$ TeV and $m_{S_+} = 1.66$ TeV, in agreement with electroweak precision tests, and plot the correlation between $\cos 2\theta$ and the product of Yukawa couplings $|y^\mu_{bL} y^\mu_{tR}|$. We scan over the values of the mixing parameter $\xi$ of Eq. (13) that are allowed by the naive perturbativity requirement, i.e., $|\xi| < \sqrt{4\pi}$. The results are illustrated in Fig. 4, where we see a very mild dependence of the product $|y^\mu_{bL} y^\mu_{tR}|$ on $\theta$ as long as we are far from $\theta \approx 0$. In other words, there is no much advantage in considering the most general case instead of the degenerate one discussed above.

4.4 Can charm work in $S_1 \& S_3$ scenario?

We briefly discuss the viability of a scenario where the leading contribution to $(g-2)_\mu$ comes not from the top quark loops but from the charm quark loops instead. In that case,
Figure 4: The product of $y_{L}^{b\mu}$ and $y_{R}^{b\mu}$ plotted against $\cos 2\theta$ for the $S_1$ & $S_3$ scenario. The masses are set to $m_{S_-} = 1.6$ TeV and $m_{S_+} = 1.66$ TeV, while the mixing parameter $\xi$, defined in Eq. (13), has been scanned over the range $|\xi| < \sqrt{4\pi}$. Dark (light) blue region is allowed by the various flavor and electroweak precision constraints described in Sec. 4.2 and addresses $(g-2)_\mu$ to 1\sigma (2\sigma) accuracy.

the chiral enhancement is not as significant as it is in the top quark case but could be, in principle, still large enough to significantly impact $\delta a_\mu$. As we are going to demonstrate, this naive reasoning is not correct due to additional constraints from flavor physics that become relevant in this scenario.

We assume that the only non-zero LQ couplings in Eqs. (24) and (25) are $y_{R}^{s\mu}$ and $y_{L}^{s\mu}$, and we focus once again on the maximal-mixing scenario. By keeping only the dominant, chirality enhanced, contributions to $\delta a_\mu$, we find that

$$\delta a_\mu \approx -2.7 \times 10^{-9} \left( \frac{\text{Re}(y_{L}^{s\mu}y_{R}^{c\mu})}{0.5} \right) \left( \frac{\delta m_{S}^{(1/3)}}{40\text{ GeV}} \right) \left( \frac{1\text{ TeV}}{m_{S}} \right)^2 ,$$

which can improve the description of $a_\mu^{\text{exp}}$ for $m_{S} \approx 1$ TeV and $O(1)$ Yukawa couplings of opposite signs, in agreement with LHC and perturbativity constraints. However, the main constraint to this scenario comes, instead, from charm physics, which precludes this possibility. Namely, LQ mixing also induces a contribution to rare charm decays, lifting the SM helicity suppression as follows,

$$B(D^{0} \rightarrow \mu\mu) = \frac{\tau_{D_{0}} f_{D}^{2} m_{D_{0}}^{3}}{64\pi} \left( \frac{\delta m_{S}^{(1/3)}}{m_{S}} \right)^2 |y_{R}^{c\mu} (V_{us} y_{L}^{s\mu})^*|^2 \left( \frac{m_{D_{0}}}{m_{c}} \right)^2 \beta_{\mu} \left( 1 - \frac{2m_{\mu}^2}{m_{D_{0}}^2} \right) ,$$

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where \( f_D = 209(3) \text{ MeV} \) [33] is \( D \)-meson decay constant, \( \beta_\mu = \sqrt{1 - 4m_\mu^2/m_{D0}^2} \) and where we have kept only the dominant LQ contribution. Current constraints tell us that \( \mathcal{B}(D^0 \rightarrow \mu\mu)^\text{exp} < 6.2 \times 10^{-9} \) [34], in such a way that one should have

\[
\left| \frac{y_L^{s\mu} y_R^{e\mu}}{m_S^2} \right| \lesssim 0.08 \text{ TeV}^{-2},
\]  

(35)

implying that the contribution in Eq. (33) can marginally improve the description of \( \delta a_\mu^\text{exp} \).

5 \( (g - 2)_\mu \) via \( \tilde{S}_1 \& S_3 \)

The \( \tilde{S}_1 \& S_3 \) scenario is rather special since LQs in question are chiral in nature. In other words, the only way they can generate chirality-enhanced contributions towards the anomalous magnetic moment of muon is if they mix through the interactions with the SM Higgs field.

The most dominant contribution towards \( (g - 2)_\mu \) in the \( \tilde{S}_1 \& S_3 \) scenario, according to Table 2, comes from bottom quark in the loops of Fig. 2. To demonstrate that, we spell out relevant interactions of \( \tilde{S}_1 \) and \( S_3 \) with the SM fermions. These are

\[
\mathcal{L}_{\tilde{S}_1} = y_{Ri}^{ij} \bar{d}^C_{Ri} e_{Rj} \tilde{S}_1 + \text{h.c.},
\]

(36)

\[
\mathcal{L}_{S_3} = y_{Li}^{ij} Q_i^C i\gamma_2 (\vec{r} \cdot \tilde{S}_3) L_j + \text{h.c.},
\]

(37)

where \( y_R \) and \( y_L \) are Yukawa coupling matrices. We omit \( B \) number violating couplings of both \( \tilde{S}_1 \) and \( S_3 \). The Lagrangians in Eqs. (36) and (37) can be expanded in terms of charge eigenstates as

\[
\mathcal{L}_{\tilde{S}_1 \& S_3} = y_{Ri}^{ij} \bar{d}^C_{Ri} e_{Rj} \tilde{S}_1^{(4/3)} - y_{Li}^{ij} \bar{Q}_i^C \nu_{Lj} S_3^{(1/3)} - \sqrt{2} y_{Li}^{ij} \bar{d}^C_{Lj} \nu_{Lj} S_3^{(1/3)}
\]

\[
+ \sqrt{2} (V^* y_L)_{ij} \bar{u}_L^C \nu_{Lj} S_3^{(2/3)} - (V^* y_L)_{ij} \bar{u}_L^C \nu_{Lj} S_3^{(1/3)} + \text{h.c.,}
\]

(38)

where we see that, in order to close the loops of Fig. 2, we need to mix the \( Q = 4/3 \) components of \( \tilde{S}_1 \) and \( S_3 \), as accomplished by the operator introduced in Eq. (14). There are four LQ mass eigenstates in this particular scenario as discussed in Sec. 3. These, again, are \( m_{S_\pm} \equiv m_{S_{4/3}^{\pm}} \) and \( m_{S_3} \equiv m_{S_{1/3}^{(1/3)}} = m_{S_{3}^{(-2/3)}} \), where we define the mass eigenstates for mixed states \( \tilde{S}_1^{(4/3)} \) and \( S_3^{(4/3)} \) as \( S_{\pm} \equiv S_{\pm}^{(4/3)} \).

We proceed following the same steps that are used to analyse the \( S_1 \& S_3 \) scenario and accordingly switch on only those Yukawa couplings in Eq. (38) that are absolutely necessary to generate contribution towards \( (g - 2)_\mu \) via bottom quark exchange. These are \( y_R^{bu} \) and \( y_L^{bu} \), where we set all other entries in \( y_R \) and \( y_L \) of Eqs. (36) and (37) to zero. The resulting analysis, in the case of the maximal mixing, i.e., when \( m_{S_3} = m_{\tilde{S}_1} \equiv m_S \) and \( \theta = \pi/4 \), yields available parameter space in the \( y_L^{bu} - y_R^{bu} \) plane that we show in Fig. 5. To produce Fig. 5 we set \( m_S = 1.6 \text{ TeV} \) and we assume that the mass splitting parameter \( \delta m_S^{(4/3)} \equiv m_S - m_{S_-} \approx m_{S_+} - m_S > 0 \) saturates electroweak precision measurement constraints with the upper limit given in Eq. (20).
Figure 5: The values of $y_{bL}^\mu$ vs. $y_{bR}^\mu$ for the $\tilde{S}_1 & S_3$ scenario that satisfy relevant flavor constraints and address $(g-2)_\mu$ are shown in dark (light) blue to 1σ (2σ) accuracy for $m_{\tilde{S}_1} = m_{S_3} = 1.6$ TeV and $\delta m_S^{4/3} = 50$ GeV. The individual 2σ constraints from $Z$-pole observables and $(g-2)_\mu$ are represented with green and red dashed lines, respectively. The (projected) LHC limits from the study of $pp \to \mu\mu$ at high-$p_T$, at the 2σ level, for $(300 \text{ fb}^{-1})$ 36 fb$^{-1}$ worth of data are shown with (dash-dotted) dashed gray lines [28].

It is clear from Fig. 5 that to accommodate $(g-2)_\mu$ within the $\tilde{S}_1 & S_3$ scenario one needs at least one of the two Yukawa couplings to be an $\mathcal{O}(1)$ parameter. Moreover, an improvement of the LHC limits from the study of $pp \to \mu\mu$ at high-$p_T$ can have significant impact on the currently available parameter space.

6 $(g-2)_\mu$ via $\tilde{R}_2 & R_2$

The dominant contribution towards $(g-2)_\mu$ in the $\tilde{R}_2 & R_2$ scenario comes from the bottom quark in the loops of Fig. 2. The relevant interactions of $\tilde{R}_2$ and $R_2$ with the SM fermions are

$$\mathcal{L}_{\tilde{R}_2} = -y_{L}^{ij} \overline{d}_{Ri}\tilde{R}_2 i\tau_2 L_j + \text{h.c.}, \quad (39)$$

$$\mathcal{L}_{R_2} = y_{R}^{ij} \overline{Q}_i e_{Rj} R_2 + \text{h.c.}. \quad (40)$$
where $y_L$ and $y_R$ are Yukawa coupling matrices. We omit the couplings of $R_2$ with right-chiral leptons to investigate the chirality-enhanced contributions coming solely from the LQ mixing scenario. The above Lagrangians can be expanded in terms of charge eigenstates as

$$\mathcal{L}_{\tilde{R}_2 \& R_2} = -y^{ij}_L \tilde{d}_R e_{Lj} \tilde{R}_2^{(2/3)} + y^{ij}_L \tilde{d}_R \nu_{Lj} \tilde{R}_2^{(-1/3)} + (V y_R)^{ij} \tilde{u}_{Li} e_{Rj} R_2^{(5/3)} + y^{ij}_R \tilde{d}_{Li} \nu_{Rj} R_2^{(2/3)} + \text{h.c.},$$

(41)

where we see that, in order to close the loops of Fig. 2, we need to mix the $Q = 2/3$ components of $\tilde{R}_2$ and $R_2$, as achieved via the operator introduced in Eq. (15). After the Higgs field gets vacuum expectation value we are left with four mass eigenstates with masses $m_{S_\pm}$, $m_{\tilde{R}_2} \equiv m_{\tilde{R}_2^{(-1/3)}}$, and $m_{R_2} \equiv m_{R_2^{(5/3)}}$, where we define the mass eigenstates for mixed states $\tilde{R}_2^{(2/3)}$ and $R_2^{(2/3)}$ as $S_{\pm} \equiv S_{\pm}^{(2/3)}$.

We consider the maximal mixing scenario only, as discussed in Sec. 3, when $m_{\tilde{R}_2} = m_{R_2} \equiv m_S$ and $\theta = \pi/4$, and, again, introduce mass splitting parameter $\delta m_s^{(2/3)} \equiv m_S - m_{S_\pm} \approx m_{S_+} - m_S > 0$. We furthermore switch on only those Yukawa couplings in Eq. (41) that are absolutely necessary to generate contribution towards $(g - 2)_\mu$ via bottom quark exchange. These are $y^{bu}_L$ and $y^{bu}_R$, where we set all other entries in $y_L$ and $y_R$ matrices, as defined in Eqs. (39) and (40), to zero. The resulting analysis yields available parameter space in the $y^{bu}_L$-$y^{bu}_R$ plane that we show in Fig. 6, where we set $m_{\tilde{R}_2} = m_{R_2} = 1.6 \, \text{TeV}$ and saturate $\delta m_s^{(2/3)}$ in agreement with the limit of Eq. (23).

It is clear from Fig. 6 that to accommodate $(g - 2)_\mu$ within the $\tilde{R}_2 \& R_2$ scenario, similarly to the $S_1$ & $S_3$ case, one needs at least one of the two Yukawa couplings to be an $\mathcal{O}(1)$ parameter. Moreover, an improvement of the LHC limits from the study of $pp \to \mu\mu$ at high-$p_T$ can have significant impact on the currently available parameter space.

7 Conclusions

We investigate viability of those scenarios where the one-loop contributions towards the anomalous magnetic moment of muon are induced through the mixing of two scalar LQs of the same electric charge $Q$ via the interactions with the SM Higgs field. The LQ pairs in question need to couple to muons and quarks of opposite chiralities in order to produce chirality-enhanced contributions. The three LQ pairs that satisfy these criteria are $S_1$ & $S_3$, $\tilde{S}_1$ & $S_3$, and $\tilde{R}_2$ & $R_2$, where the two states that mix have the electric charges $Q = 1/3$, $Q = 4/3$, and $Q = 2/3$, respectively. Since, in all three instances, at least one of the LQ multiplets carries non-trivial $SU(2)_L$ assignment the mixing can induce mass splitting between the states belonging to the same LQ multiplet that, in turn, might generate substantial oblique corrections that could potentially be in conflict with the existing electroweak precision measurements. The electroweak observable $T$ imposes the constraint on the mass splitting that we, in the case of maximal mixing between the two LQs, describe with a parameter $\delta m_s^{(0)}$ to find $\delta m_s^{(1/3)} \lesssim 40 \, \text{GeV}$, $\delta m_s^{(4/3)} \lesssim 50 \, \text{GeV}$, and $\delta m_s^{(2/3)} \lesssim 50 \, \text{GeV}$ for the $S_1$ & $S_3$, $\tilde{S}_1$ & $S_3$, and $\tilde{R}_2$ & $R_2$ scenarios, respectively.
Figure 6: The values of $y_{R}^{b\mu}$ vs. $y_{L}^{b\mu}$ for the $R_{2}$ & $R_{2}$ scenario that satisfy relevant flavor constraints and address $(g-2)_{\mu}$ are shown in dark (light) blue to $1\sigma$ ($2\sigma$) accuracy for $m_{R_{2}} = m_{R_{2}} = 1.6\text{ TeV}$ and $\delta m_{S}^{(2/3)} = 50\text{ GeV}$. The individual $2\sigma$ constraints from $Z$-pole observables and $(g-2)_{\mu}$ are represented with green and red dashed lines, respectively. The (projected) LHC limits from the study of $pp \to \mu\mu$ at high-$p_{T}$, at the $2\sigma$ level, for $(300\text{ fb}^{-1})$ 36 fb$^{-1}$ worth of data are shown with (dash-dotted) dashed gray lines [28].

We considered, in all three instances, the most minimal set of Yukawa couplings that can generate chirality-enhanced contribution to $(g-2)_{\mu}$ at the one-loop level. This contribution is proportional to the mass splitting parameter $\delta m_{S}^{(Q)}$ and the product of one left- and one right-chiral coupling of LQs to muons whereas it is inversely proportional to the common scale for the LQ masses that we set at 1.6 TeV for our numerical analysis.

In the $S_{1}$ & $S_{3}$ case the quarks in the loop are the up-type ones and we investigate viability of both the top quark and the charm quark cases, where, in addition to the electroweak precision measurement constraints, we impose relevant input from flavor physics, in particular $Z \to lll$, and the current LHC data analyses. We find that the top-quark loop contributions represent a viable option whereas the charm-quark loop contributions turn out to be in direct conflict with the flavor physics constraints. If we allow, in the top quark case, that the relevant Yukawa couplings can reach the perturbativity limit, we find that the relevant scale for the LQ masses should be smaller than 15 TeV if one is to address $(g-2)_{\mu}$.

In the $\tilde{S}_{1}$ & $S_{3}$ and $\tilde{R}_{2}$ & $R_{2}$ scenarios the quarks in the $(g-2)_{\mu}$ loops are the down-type
ones and we accordingly consider the bottom quark contributions only. We find, after imposing all relevant flavor physics and LHC constraints, that both scenarios require at least one of the two Yukawa couplings to be an $\mathcal{O}(1)$ parameter to be able to address the $(g - 2)_\mu$ discrepancy. The existing LHC limits from the study of $pp \to \mu\mu$ at high-$p_T$ somewhat constrain available parameter space of these two scenarios while the projected limits that correspond to the $300\,\text{fb}^{-1}$ worth of the LHC data would significantly impair their ability to accommodate $(g - 2)_\mu$.

Even though we mostly consider the limit when the two LQs mix maximally we also demonstrate that the departure from that limit is not more beneficial for the ability of these three scenarios to address the anomalous magnetic moment of muon discrepancy.

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