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Neutrino mass, mixing and muon $g-2$ explanation in $\mathrm{U}(1)_{L_\mu - L_\tau}$ extension of left-right theory

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Abstract: We consider a gauged $\mathrm{U}(1)_{L_\mu - L_\tau}$ extension of the left-right symmetric theory in order to simultaneously explain neutrino mass, mixing and the muon anomalous magnetic moment. We get sizeable contribution from the interaction of the new light gauge boson $Z_{\mu\tau}$ of the $\mathrm{U}(1)_{L_\mu - L_\tau}$ symmetry with muons which can individually satisfy the current bounds on muon $(g-2)$ anomaly (Δa_μ). The other positive contributions to Δa_μ come from the interactions of singly charged gauge bosons W_L, W_R with heavy neutral fermions and that of neutral CP-even scalars with muons. The interaction of W_L with heavy neutrino is facilitated by inverse seesaw mechanism which allows large light-heavy neutrino mixing and explains neutrino mass in our model. CP-even scalars with mass around few hundreds GeV can also satisfy the entire current muon anomaly bound. The results show that the model gives a small but non-negligible contribution to Δa_μ thereby eliminating the entire deviation in theoretical prediction and experimental result of muon $(g-2)$ anomaly. We have briefly presented a comparative study for symmetric and asymmetric left-right symmetric model in context of various contribution to Δa_μ. We also discuss how the generation of neutrino mass is affected when left-right symmetry breaks down to Standard Model symmetry via various choices of scalars.

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JHEP09(2020)010

Contents

1 Introduction

While most of the theoretical predictions by Standard Model (SM) have been experimentally found to be correct to a very high precision, there lies a wide gap between SM's prediction of muon anomalous magnetic moment, $a_\mu = \frac{g_\mu - 2}{2}$ and its measurement. The SM prediction can be summed up as $a_\mu^{\rm SM} = (11659183.0 \pm 4.8) \times 10^{-10}$ [1, 2] whereas, the value obtained by Brookhaven National Laboratory (BNL) is $a_\mu^{\rm exp} = (11659209.1 \pm 6.3) \times 10^{-10}$ [1, 3] with $\Delta a_\mu = (26.1 \pm 7.9) \times 10^{-10}$ [4]. While a 3.3σ deviation is achieved by BNL yet [3], a nearly 5σ deviation is expected in the near future by Fermilab E989 [5] and of similar precision by J-PARC [6]. In principle the a_μ predicted by SM is a sum of contributions coming from QED, electroweak and hadronic sectors;

$$a_\mu^{\rm SM} = a_\mu^{\rm QED} + a_\mu^{\rm electroweak} + a_\mu^{\rm hadronic} \tag{1.1}$$

Among these three contributions, the theoretical uncertainty is believed to be coming from the hadronic loop contributions [7–9] since the other two contributions have been verified with a high precision [10, 11]. A proposed experiment, namely MUonE [12] aspires to reduce this theoretical uncertainty by determining the hadronic vacuum polarization more precisely. All these recent developments in the experimental muon sector surely ignites theoretical research that aim at eliminating or narrowing down this wide gap in the prediction and measurement. Therefore recently many new physics scenarios have been explored in this context, for an incomplete list of which one may refer [13–34].

Many of these new physics scenarios focus on $\mathrm{U}(1)_{L_\mu - L_\tau}$ symmetry to address the anomaly because of the phenomenology associated with its gauge boson $Z_{\mu\tau}$. The total lepton number, L, is a sum of individual lepton numbers L_e, L_μ, L_τ and one can always choose the difference between any two individual lepton numbers like $L_e - L_\mu$, $L_\mu - L_\tau$, $L_e - L_\tau$ and gauge it to obtain an anomaly free theory. However, the gauged $\mathrm{U}(1)_{L_\mu - L_\tau}$ symmetry is the most chosen one due to the fact that the parameters associated with $Z_{\mu\tau}$ gauge boson is not constrained by lepton and hadron colliders since it doesn't couple to electrons and quarks. Moreover, as per the constraints given by neutrino-trident experiments [35] a low mass of $\mathcal{O}(100\,\mathrm{MeV})$ can be allowed for this new gauge boson $Z_{\mu\tau}$ for a coupling as low as $g_{\mu\tau} \leq 10^{-3}$.

The $\mathrm{U}(1)_{L_\mu - L_\tau}$ extension of SM has been extensively studied for explaining several issues like muon $(g-2)$ anomaly [36, 37], dark matter [38], orbital energy loss of a neutron star [39] and so on. Several other works have explained how the associated $Z_{\mu\tau}$ gauge boson can ameliorate the tension in the late time and early time determination of Hubble constant [40], unexpected dip in the energy spectrum of high energy cosmic neutrinos reported by the IceCube Collaboration [41] and also the deviations to neutrino oscillations due to long range forces [42]. Ref. [43] says the vectors associated with a gauged $\mathrm{U}(1)_{L_\mu - L_\tau}$ symmetry can induce an anomalously fast decay of the orbital period of neutron star binaries which might be used to discover any long-ranged muonic force associated with the binaries while ref [36] explains how this gauge boson can possibly mediate interactions between dark matter particles and muons inside a neutron star. The effect induced by $L_\mu - L_\tau$ vector to enhanced production in neutrino decays, meson decays, neutrinoless double beta decays, and annihilations are discussed in ref [44]. The possible detection of this light $Z_{\mu\tau}$ boson is discussed in refs. [45–48]. However the $\mathrm{U}(1)_{L_\mu - L_\tau}$ extension of SM can not accommodate neutrino mass until and unless one adds a right-handed neutrino to the model. Such attempts have been made in ref [38, 49] , where the authors explain neutrino mass by adding three right-handed neutrinos to the model. In ref [50] neutrino masses with bimaximal mixing is obtained just by adding one right-handed neutrino to the extended SM framework. A similar framework [51] also predicts quasi-degenerate neutrino masses. On the other hand, the left-right symmetric model (LRSM) [52–59] is a SM extension which clearly gives us a unified answer to small neutrino mass generation as well as parity violation problem in low-energy weak interactions. LRSM naturally hosts a right-handed neutrino and offers wider possibilities of explaining neutrino mass, lepton number violation, lepton flavour violation with rich phenomenology at low scale. In particular we shall see that an interplay between the right handed gauge bosons and the mass mechanisms for the neutrinos makes an important contribution.

Thus with the motivation of explaining neutrino mass, mixing and muon $(g-2)$ anomaly in a single framework we reach for the LRSM and augment it with the $\mathrm{U}(1)_{L_\mu - L_\tau}$ symmetry. In manifest LRSM neutrino mass can be explained by canonical seesaw mechanism, but it cannot be verified by collider experiments since a very high right-handed breaking scale ($10^{14}\,\mathrm{GeV}$) is associated with the mechanism. Thus in general extra particles are added to LRSM in order to generate neutrino mass by various low-scale seesaw mechanisms like linear seesaw, inverse seesaw [60–70], double seesaw etc [68, 71–83]. In particular, we take interest in inverse seesaw in our extended LRSM to explain neutrino mass which also allows large light-heavy neutrino mixing and thus leads to sizeable contributions to the muon anomalous magnetic moment via left-handed singly-charged SM gauge boson interaction with heavy neutrino. Apart from the usual fermions and scalars present in a manifest LRSM, the model contains three sets of extra sterile fermions and one extra scalar. While the extra sterile fermions help in creating the plot for inverse seesaw, the extra scalar helps in breaking the $\mathrm{U}(1)_{L_\mu - L_\tau}$ symmetry and also in implementing the inverse seesaw in the model. The $Z_{\mu\tau}$ boson originated from the breaking of $\mathrm{U}(1)_{L_\mu - L_\tau}$ symmetry helps in ameliorating Δa_μ when it gets mass around 150 MeV. Moreover our predictions on the mass of $Z_{\mu\tau}$ and its coupling $g_{\mu\tau}$ lie well below the constraint given by ref [84]. We also discuss various symmetry breaking chains from LRSM to SM with different choices of scalars to see how it affects the generation of neutrino mass. Also we have shown that lighter neutral CP-even scalars can also satisfy the current as well as 1σ bound on muon anomaly individually if they possess mass around 0.5–2 TeV.

The rest of the paper is organised as follows. In section 2 we present the particle content of the extended LRSM and discuss the symmetry breaking of $\mathrm{U}(1)_{L_\mu - L_\tau}$ symmetry and left-right symmetry down to low energy theory. We also discuss two different scenarios of neutrino mass generation with the help of doublet scalars in 2.1 and triplet scalars in 2.2. In section 3 we discuss the generation of neutrino mass and mixing via extended inverse seesaw mechanism. In section 4 we analytically study the new contributions to Δa_μ arising from different vector bosons and scalars present in the model. In section 5 we estimate the contributions numerically and present the results. This section also contains several plots of Δa_μ *vs* mass of mediators to check the sensitivity of our theoretical results to experimental bounds. In section 6 we summarize and conclude the work.

2 The model

The model is an extension of manifest left-right theory with additional U(1) gauge symmetry where the difference between muon and tau lepton numbers is gauged. The model is governed by the gauge group,

$$\mathbb{G}_{\mathbb{LR}}^{\mu\tau} \equiv \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L} \times \mathrm{SU}(3)_C \times \mathrm{U}(1)_{L_\mu - L_\tau} \tag{2.1}$$

Within manifest LRSM which is based on the gauge group $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L} \times \mathrm{SU}(3)_C$ and consists of usual quarks ($q_{L,R}$), leptons ($\ell_{L,R}$), Higgs bidoublet Φ and triplets $\Delta_{L,R}$ (presented in table 1) the light neutrino masses can be generated by type-I+II seesaw mechanism [72, 73, 76, 77, 80, 85, 86],

$$m_\nu = -M_D M_R^{-1} M_D^T + M_L = m_\nu^I + m_\nu^{II} \,,$$

where $M_L(M_R)$ represents the Majorana mass term for light left-handed (heavy right-handed) Majorana neutrinos arising from respective VEVs of left-handed (right-handed) scalar triplet and M_D is the Dirac neutrino mass matrix connecting light-heavy neutrinos. Here, the scale of right-handed neutrino mass (M_R) is related to the non-zero VEV of right-handed scalar triplet which is responsible for spontaneous symmetry breaking of LRSM

JHEP09(2020)010

	Fields	$SU(2)_L$	$SU(2)_R$	$B-L$	$SU(3)_C$
Fermions	q_L	2	1	1/3	3
	q_R	1	2	1/3	3
	ℓ_L	2	1	-1	1
	ℓ_R	1	2	-1	1
Scalars	Φ	2	2	0	1
	Δ_L	3	1	2	1
	Δ_R	1	3	2	1

Table 1. Particle content of the manifest left-right symmetric theories.

to SM. The sub-eV scale of light neutrino mass, as hinted by oscillation experiments, is connected to a very heavy right-handed scale i.e, 10^{15} GeV (in generic scenarios) clearly making it inaccessible to current and planned accelerator experiments. On the other hand, when LRSM breaks around TeV scale, the gauge bosons W_R, Z_R, right-handed neutrinos N_R and scalar triplets $\Delta_{L,R}$ get TeV scale mass that allows several lepton number violating signatures at LHC as well as low energy experiments like neutrinoless double beta decay. The left-right mixing (or light-heavy neutrino mixing), which depends on Dirac neutrino mass M_D, plays an important role in giving large new contribution to neutrinoless double beta decay, other LNV signatures at colliders as well as LFV processes. This gives the motivation to explore alternative class of left-right symmetric model with large value of M_D and thereby large light-heavy neutrino mixing which can contribute positively to Δa_μ.

A number of LRSM variants have been explored in literature [81, 87–92] where spontaneous symmetry breaking is implemented with scalar bidoublet having $B-L=0$ and Higgs doublets having $B-L=1$ which leads to neutrino mass being generated by either simple Dirac mass terms or low scale seesaw mechanisms like inverse seesaw, linear seesaw etc. In this model, for the generation of neutrino mass we take interest in inverse seesaw mechanism since it allows large light-heavy neutrino mixing and this mixing facilitates the interaction of singly charged vector boson with heavy neutrinos which contributes positively to Δa_μ. Before we move on to the working of inverse seesaw mechanism in the considered model, let's have a clear picture of how the generation of neutrino mass is affected within various symmetry breaking of LRSM-SM chains.

At first, the spontaneous symmetry breaking (SSB) of $\mathbb{G}_{\mathbb{LR}}^{\mu\tau}$ down to left-right theory $\mathbb{G}_{\mathbb{LR}}$ is achieved by assigning a non-zero VEV to a scalar χ which is singlet under left-right symmetry but non-trivially charged under $U(1)_{L_\mu - L_\tau}$. Further, the SSB of LRSM to SM can happen in the following three ways;

- with Higgs doublets $H_L \oplus H_R$,
- with Higgs triplets $\Delta_L \oplus \Delta_R$,
- with the combination of doublets and triplets $H_L \oplus H_R$ and $\Delta_L \oplus \Delta_R$.

Now, as usual the SSB of SM to low energy theory occurs when the scalar bidoublet Φ takes non-zero vev and that generates masses for charged leptons and quarks.

JHEP09(2020)010

	Fields	$SU(2)_L$	$SU(2)_R$	$U(1)_{B-L}$	$SU(3)_C$	$U(1)_{L_\mu-L_\tau}$
Fermions	ℓ_{e_L}	2	1	-1	1	0
	ℓ_{μ_L}	2	1	-1	1	1
	ℓ_{τ_L}	2	1	-1	1	-1
	ℓ_{e_R}	1	2	-1	1	0
	ℓ_{μ_R}	1	2	-1	1	1
	ℓ_{τ_R}	1	2	-1	1	-1
Scalars	Φ	2	2	0	1	0
	H_L	2	1	1	1	0
	H_R	1	2	1	1	0
	χ	1	1	0	1	1 or 2

Table 2. Particle content of left-right theories extended with $U(1)_{L_\mu-L_\tau}$ gauge symmetry where fermion sector is limited to leptons and scalar sector contains the bidoublet Φ, doublets $H_{L,R}$ and a singlet χ.

2.1 Neutrino masses with LRSM-SM symmetry breaking via H_R, H_L

In this minimal version, H_R breaks the left-right symmetry to SM while H_L is required for left-right invariance. The scalar bidoublet Φ is required for SM symmetry breaking to low energy theory and χ is needed for the spontaneous symmetry breaking (SSB) of $\mathbb{G}_{\mathbb{LR}}^{\mu\tau}$ down to left-right LR theory $\mathbb{G}_{\mathbb{LR}}$ as mentioned earlier. The leptons and scalars are displayed in table 2. The allowed Yukawa interactions for leptons are given by,

$$-\mathcal{L}_{Yuk} \supset \overline{\ell_{e_L}}\left[Y_\ell\Phi + \tilde{Y}_\ell\widetilde{\Phi}\right]\ell_{e_R} + \overline{\ell_{\mu_L}}\left[Y_\ell\Phi + \tilde{Y}_\ell\widetilde{\Phi}\right]\ell_{\mu_R} + \overline{\ell_{\tau_L}}\left[Y_\ell\Phi + \tilde{Y}_\ell\widetilde{\Phi}\right]\ell_{\tau_R} + \text{h.c.} \tag{2.2}$$

with $\tilde{\Phi} = \tau_2\Phi^*\tau_2$.

The vev structure for the Higgs spectrum can be depicted as follows:

$$\langle H_R\rangle = \begin{pmatrix} 0 \\ v_R \end{pmatrix}, \quad \langle H_L\rangle = \begin{pmatrix} 0 \\ v_L \end{pmatrix}, \quad \langle\Phi\rangle = \begin{pmatrix} v_1 & 0 \\ 0 & v_2 \end{pmatrix}.$$

After SSB, the charged fermion as well as light neutrino mass matrices are found to be diagonal in structure due to presence of $U(1)_{L_\mu-L_\tau}$ gauge symmetry. This is the important prediction of the model giving simplified relation for PMNS mixing matrix as $U_{\text{PMNS}} \equiv U_\nu$.

The non-zero masses for light neutrinos (which are Dirac fermions) can be explained by adjusting Yukawa couplings through the non-zero VEVs of scalar bidoublet. From the Yukawa interactions given in eq. (2.2); with $Y_\ell \ll \tilde{Y}_\ell$, $v_2 \ll v_1$, the masses for charged leptons and the light neutrinos can be expressed as,

$$M_\ell \simeq \tilde{Y}_\ell v_1^*\,, \qquad M_D^\nu \simeq v_1\left(Y_\ell + M_\ell\frac{v_2}{v_1^2}\right). \tag{2.3}$$

Even though this framework holds a minimal (in terms of SU(2) representation) scalar spectrum it can not provide Majorana mass for neutrinos and thus forbids any signature of lepton number violation.

2.2 Neutrino masses with LRSM-SM symmetry breaking via Δ_R, Δ_L

In table 2 if we replace the doublets H_L, H_R by triplets Δ_L, Δ_R then the model offers a better possibility from phenomenology point of view since in this case Majorana masses can be generated for light and heavy neutrinos. If the symmetry breaking occurs at few TeV scale, these Majorana neutrinos can mediate neutrinoless double beta decay process whose observation would confirm lepton number violation in nature. Lepton number violation can also be probed via smoking-gun same-sign dilepton signatures at collider experiments. The interaction terms involving scalar triplets and leptons in the left-right theories with extra U(1) symmetry are given by

$$\begin{aligned}
-\mathcal{L}_{Yuk} \supset\ & \overline{\ell_{e_L}} \left[Y_\ell \Phi + \tilde{Y}_\ell \widetilde{\Phi}\right] \ell_{e_R} + \overline{\ell_{\mu_L}} \left[Y_\ell \Phi + \tilde{Y}_\ell \widetilde{\Phi}\right] \ell_{\mu_R} + \overline{\ell_{\tau_L}} \left[Y_\ell \Phi + \tilde{Y}_\ell \widetilde{\Phi}\right] \ell_{\tau_R} \\
& + \left[f_{ee} \overline{(\ell_{e_L})^c} \ell_{e_L} + f_{\mu\tau} \overline{(\ell_{\mu_L})^c} \ell_{\tau_L} + f_{\tau\mu} \overline{(\ell_{\tau_L})^c} \ell_{\mu_L} \right] \Delta_L \\
& + \left[f_{ee} \overline{(\ell_{e_R})^c} \ell_{e_R} + f_{\mu\tau} \overline{(\ell_{\mu_R})^c} \ell_{\tau_R} + f_{\tau\mu} \overline{(\ell_{\tau_R})^c} \ell_{\mu_R} \right] \Delta_R + \text{h.c.}
\end{aligned} \tag{2.4}$$

with the corresponding vevs

$$\langle \Delta_L \rangle = \begin{pmatrix} 0 & 0 \\ v_L & 0 \end{pmatrix}, \quad \langle \Delta_R \rangle = \begin{pmatrix} 0 & 0 \\ v_R & 0 \end{pmatrix}, \quad \langle \Phi \rangle = \begin{pmatrix} v_1 & 0 \\ 0 & v_2 \end{pmatrix}$$

Using eq. (2.4), the structure of the masses for neutral leptons in the basis (ν_L, N_R^c) can be written as,

$$\mathbb{M} = \begin{pmatrix} M_L & M_D \\ M_D^T & M_R \end{pmatrix}, \tag{2.5}$$

where, M_D represents Dirac neutrino mass matrix, $M_L (M_R)$ denotes Majorana mass matrix arising from the non-zero vev of LH (RH) scalar triplet. The mass matrices M_D, M_L and M_R can be written explicitly as follows (considering $f_{\mu\tau} = f_{\tau\mu}$ and $f_{\mu\mu} = f_{\tau\tau}$ for sake of simplicity),

$$\begin{aligned}
M_D &= \begin{pmatrix} Y_{11} v_2 + \tilde{Y}_{11} v_1 & 0 & 0 \\ 0 & Y_{22} v_2 + \tilde{Y}_{22} v_1 & 0 \\ 0 & 0 & Y_{33} v_2 + \tilde{Y}_{33} v_1 \end{pmatrix} = \begin{pmatrix} a & 0 & 0 \\ 0 & b & 0 \\ 0 & 0 & c \end{pmatrix}, \\
M_{L,R} &= \begin{pmatrix} f_{ee} & 0 & 0 \\ 0 & 0 & f_{\mu\tau} \\ 0 & f_{\mu\tau} & 0 \end{pmatrix} \frac{v_{L,R}}{\sqrt{2}},
\end{aligned} \tag{2.6}$$

Now using seesaw approximation $M_R \gg M_D$ and $M_L \to 0$, the light neutrino mass can be generated via type-I seesaw formula as shown below,

$$\begin{aligned}
m_\nu^I &= -M_D\, M_R^{-1}\, M_D^T \\
&= \begin{pmatrix} a & 0 & 0 \\ 0 & b & 0 \\ 0 & 0 & c \end{pmatrix} \cdot \begin{pmatrix} f_{ee} \frac{v_R}{\sqrt{2}} & 0 & 0 \\ 0 & 0 & f_{\mu\tau} \frac{v_R}{\sqrt{2}} \\ 0 & f_{\mu\tau} \frac{v_R}{\sqrt{2}} & 0 \end{pmatrix}^{-1} \cdot \begin{pmatrix} a & 0 & 0 \\ 0 & b & 0 \\ 0 & 0 & c \end{pmatrix}^T = \begin{pmatrix} \frac{\sqrt{2} a^2}{f_{ee} v_R} & 0 & 0 \\ 0 & 0 & \frac{\sqrt{2} bc}{f_{\mu\tau} v_R} \\ 0 & \frac{\sqrt{2} bc}{f_{\mu\tau} v_R} & 0 \end{pmatrix}
\end{aligned} \tag{2.7}$$

From this light neutrino mass matrix m_ν^I, the corresponding mass eigenvalues for light neutrino mass eigenstates can be obtained which are, $\left\{-\frac{\sqrt{2}bc}{f_{\mu\tau}v_R}, \frac{\sqrt{2}bc}{f_{\mu\tau}v_R}, \frac{\sqrt{2}a^2}{f_{ee}v_R}\right\}$. However, two mass eigenstates with eigenvalues $\frac{\sqrt{2}bc}{f_{\mu\tau}v_R}$ (ignoring the relative negative sign) are degenerate here, which implies either the solar neutrino mass difference ($\Delta m^2_{\rm sol}$) or the atmospheric neutrino mass difference ($\Delta m^2_{\rm atm}$) vanishes. This is in disagreement with the neutrino experimental data at 3σ interval of global fit by NuFIT 4.1 [93],

$$\begin{aligned} \text{NO}: \Delta m^2_{\rm atm} &= [2.432, 2.618] \times 10^{-3}\ \text{eV}^2,\ \Delta m^2_{\rm sol} = [6.79, 8.01] \times 10^{-5}\ \text{eV}^2, \qquad (2.8)\\ \sin^2\theta_{13} &= [0.02046, 0.02440],\ \sin^2\theta_{23} = [0.427, 0.609],\ \sin^2\theta_{12} = [0.275, 0.350],\\ \text{IO}: \Delta m^2_{\rm atm} &= [2.416, 2.603] \times 10^{-3}\ \text{eV}^2,\ \Delta m^2_{\rm sol} = [6.79, 8.01] \times 10^{-5}\ \text{eV}^2, \qquad (2.9)\\ \sin^2\theta_{13} &= [0.02066, 0.02461],\ \sin^2\theta_{23} = [0.430, 0.612],\ \sin^2\theta_{12} = [0.275, 0.350]. \end{aligned}$$

This degeneracy can be wiped out by introducing another pair of triplet scalars $\Delta'_L \oplus \Delta'_R$ with $L_\mu - L_\tau = 2$. Now we can write additional Yukawa terms allowed by the $\mathrm{U}(1)_{L_\mu - L_\tau}$ symmetry as,

$$-\mathcal{L}^{\rm new}_{Yuk} \supset f_{\mu\mu}(\ell^T_{\mu R}\Delta'^\dagger_R \ell_{\mu R} + \ell^T_{\tau R}\Delta'_R \ell_{\tau R}) + R \leftrightarrow L \tag{2.10}$$

With these new permissible terms in the Yukawa sector, we can write the corresponding $M'_{L,R}$ matrix as,

$$M'_{L,R} = \begin{pmatrix} f_{ee}\frac{v_{L,R}}{\sqrt{2}} & 0 & 0 \\ 0 & f_{\mu\mu}\frac{v'_{L,R}}{\sqrt{2}} & f_{\mu\tau}\frac{v_{L,R}}{\sqrt{2}} \\ 0 & f_{\mu\tau}\frac{v_{L,R}}{\sqrt{2}} & f_{\mu\mu}\frac{v'_{L,R}}{\sqrt{2}} \end{pmatrix} \tag{2.11}$$

where $v'_{L,R} = \langle \Delta'_{L,R} \rangle$. Now using the seesaw approximation $M'_R \gg M_D$ and $M'_L \to 0$, the light neutrino mass matrix can be expressed via type-I seesaw formula as,

$$\begin{aligned} m'^I_\nu &= -M_D M'^{-1}_R M^T_D \\ &= \begin{pmatrix} \frac{\sqrt{2}a^2}{f_{ee}v_R} & 0 & 0 \\ 0 & \frac{\sqrt{2}b^2 f_{\mu\mu}v'_R}{-f^2_{\mu\tau}v^2_R + f^2_{\mu\mu}v'^2_R} & \frac{\sqrt{2}bcf_{\mu\tau}v_R}{f^2_{\mu\tau}v^2_R - f^2_{\mu\mu}v'^2_R} \\ 0 & \frac{\sqrt{2}bcf_{\mu\tau}v_R}{f^2_{\mu\tau}v^2_R - f^2_{\mu\mu}v'^2_R} & \frac{\sqrt{2}c^2 f_{\mu\mu}v'_R}{-f^2_{\mu\tau}v^2_R + f^2_{\mu\mu}v'^2_R} \end{pmatrix} \end{aligned} \tag{2.12}$$

Here all three mass eigenvalues are non-degenerate which can be represented as, $\left\{\frac{\sqrt{2}a^2}{f_{ee}v_R}, m^{Ia}_\nu \pm m^{Ib}_\nu\right\}$ with

$$\begin{aligned} m^{Ia}_\nu &= \frac{-b^2 f_{\mu\mu}v'_R - c^2 f_{\mu\mu}v'_R}{\sqrt{2}(f^2_{\mu\tau}v^2_R - f^2_{\mu\mu}v'^2_R)}, \\ m^{Ib}_\nu &= \frac{\sqrt{4b^2c^2 f^2_{\mu\tau}v^2_R + b^4 f^2_{\mu\mu}v^2_R - 2b^2c^2 f^2_{\mu\mu}v'^2_R + c^4 f^2_{\mu\mu}v'^2_R}}{\sqrt{2}(f^2_{\mu\tau}v^2_R - f^2_{\mu\mu}v'^2_R)} \end{aligned}$$

Though the introduction of two extra scalar triplets $\Delta'_L \oplus \Delta'_R$ saves us from apparent inconsistency in the explanation of current-day neutrino oscillation data, the particle content of the model becomes crowded and it no more remains minimal.

3 LRSM Inverse Seesaw (LISS) for neutrino masses

We have already discussed in previous section that the sub-eV scale neutrino masses can be generated either by canonical see-saw mechanism which requires a very high ($> 10^{14}$ GeV) seesaw scale and therefore cannot be verified by present-day colliders or with very much suppressed value of Dirac neutrino Yukawa coupling. As an alternative, we explain one of the low scale seesaw mechanism, i.e, LRSM inverse seesaw (LISS) [60, 94–105] in our model where the left-right symmetry breaking occurs at few TeV. This symmetry breaking generates TeV scale masses for W_R, Z_R gauge bosons which fall within the LHC range and the inverse seesaw mechanism provides large light-heavy neutrino mixing. As a result the large mixing between sub-TeV scale heavy neutrinos with sub-eV scale light neutrinos, within left-right inverse seesaw scheme, offers,

- sizeable contribution to muon $g-2$ anomaly arising form purely left-handed currents with the exchange of sub-TeV masses for sterile neutrinos in LISS scheme,
- dominant contribution to lepton flavour violating (LFV) decays, non-unitarity effects in leptonic sector,
- interesting collider signatures verifiable at LHC.

For implementing LISS, we consider an extra sterile neutrino S_L per generation along with the usual leptons, scalars (bidoublet Φ, doublets $H_{L,R}$ and χ) presented in table 2. The relevant Yukawa interaction Lagrangian for LISS invariant under $\mathrm{U}(1)_{L_\mu - L_\tau}$ symmetry is given as sum of different components,

$$-\mathcal{L}_{\rm LISS} = \mathcal{L}_{\nu_L N_R} + \mathcal{L}_{N_R S_L} + \mathcal{L}_{S_L S_L}\,, \tag{3.1}$$

where the individual components are given as follows:

Generic Dirac neutrino mass matrix, $\mathcal{L}_{\nu_L N_R}$: the usual Dirac Yukawa interaction Lagrangian that allows Dirac mass terms for charged leptons and neutrinos consistent with the $\mathrm{U}(1)_{L_\mu - L_\tau}$ gauge symmetry is given by,

$$\begin{aligned}\mathcal{L}_{\nu_L N_R} &\supset \overline{\ell_L}(Y\Phi + \tilde{Y}\tilde{\Phi})\ell_R \\ &= \overline{\ell_{e_L}}[M_i]^{ee}\ell_{e_R} + \overline{\ell_{\mu_L}}[M_i]^{\mu\mu}\ell_{\mu_R} + \overline{\ell_{\tau_L}}[M_i]^{\tau\tau}\ell_{\tau_R}\end{aligned} \tag{3.2}$$

where, $M_i = M_\ell, M_D^\nu \equiv M_D$ are the corresponding Dirac mass matrices for charged leptons and neutrinos respectively. The imposition of extra $\mathrm{U}(1)_{L_\mu - L_\tau}$ symmetry to the left-right theories results in diagonal Dirac mass matrices for charged leptons and neutrinos as,

$$\begin{aligned}M_\ell &= \begin{pmatrix} Y_{11}v_1 + \tilde{Y}_{11}v_2 & 0 & 0 \\ 0 & Y_{22}v_1 + \tilde{Y}_{22}v_2 & 0 \\ 0 & 0 & Y_{33}v_1 + \tilde{Y}_{33}v_2 \end{pmatrix} \\ M_D &= \begin{pmatrix} Y_{11}v_2 + \tilde{Y}_{11}v_1 & 0 & 0 \\ 0 & Y_{22}v_2 + \tilde{Y}_{22}v_1 & 0 \\ 0 & 0 & Y_{33}v_2 + \tilde{Y}_{33}v_1 \end{pmatrix} = \begin{pmatrix} a & 0 & 0 \\ 0 & b & 0 \\ 0 & 0 & c \end{pmatrix}.\end{aligned} \tag{3.3}$$

JHEP09(2020)010

Dirac Mass term between N_R and S_L, $\mathcal{L}_{N_R S_L}$: the corresponding Yukawa term gives rise to the mixing matrix M between N_R and S_L as,

$$\mathcal{L}_{N_R S_L} \supset Y_{RS} \overline{\ell} \tilde{H}_R S_L = Y_{RS} \langle \tilde{H}_R \rangle \left[\overline{\ell_{eR}} S_{e_L} + \overline{\ell_{\mu R}} S_{\mu_L} + \overline{\ell_{\tau R}} S_{\tau_L} \right] \tag{3.4}$$

The corresponding mixing matrix is also found to be diagonal as,

$$M = \begin{pmatrix} M_{11} & 0 & 0 \\ 0 & M_{22} & 0 \\ 0 & 0 & M_{33} \end{pmatrix}$$

whose diagonal entries are proportional to $\langle \tilde{H}_R \rangle = v_R$.

Bare Majorana mass term for S_L, $\mathcal{L}_{S_L S_L}$: now, we focus on the generation of bare Majorana mass term for sterile neutrinos and the $\mathrm{U}(1)_{L_\mu - L_\tau}$ gauge group allows the terms involving extra sterile neutrinos as,

$$\begin{aligned} \mathcal{L}_{S_L S_L} &= \mu S_L^T S_L \\ &= \left[\mu_{ee} S_{e_L}^T S_{e_L} + \mu_{\mu\tau} S_{\mu_L}^T S_{\tau_L} + \mu_{\mu\tau} S_{\tau_L}^T S_{\mu_L} \right] \end{aligned} \tag{3.5}$$

So the bare Majorana mass matrix structure for extra sterile neutrinos can be expressed as,

$$\mu = \begin{pmatrix} \mu_{ee} & 0 & 0 \\ 0 & 0 & \mu_{\mu\tau} \\ 0 & \mu_{\mu\tau} & 0 \end{pmatrix} \tag{3.6}$$

Thus, the complete 9×9 neutral fermion mass matrix in the basis of (ν_L, N_R, S_L) is read as,

$$\mathbb{M} = \begin{pmatrix} 0 & M_D & 0 \\ M_D^T & 0 & M^T \\ 0 & M & \mu \end{pmatrix} \tag{3.7}$$

Using eq. (3.7) with mass hierarchy $M > M_D \gg \mu$, we can write the expression for Majorana mass (m_ν) for light neutrinos and pseudo-Dirac mass term (m_H) for heavy neutrinos in LISS as [34, 95–105],

$$m_\nu = \left(\frac{M_D}{M} \right) \mu \left(\frac{M_D}{M} \right)^T \tag{3.8}$$

$$m_H = -\left(\pm M - \mu/2 \right) \tag{3.9}$$

The beautiful aspect of the low scale inverse seesaw scheme is that it allows sub-eV scale of light neutrinos with large value of M_D and M as,

$$\left(\frac{m_\nu}{0.1\,\mathrm{eV}} \right) = \left(\frac{M_D}{100\,\mathrm{GeV}} \right)^2 \left(\frac{\mu}{\mathrm{keV}} \right) \left(\frac{M}{10^4\,\mathrm{GeV}} \right)^{-2} .$$

Even with $M(\sim$ sub TeV scale), we can have sizeable light-heavy neutrino mixing $(M_D/M \simeq \mathcal{O}(0.1-1))$ which can give rise to large charged LFV decay channels as

JHEP09(2020)010

$\mu \to e\gamma, \tau \to \mu$ and $0\nu\beta\beta$ effects [98]. Now from eq. (3.8), in this LRSM inverse seesaw approximation, we can express the light neutrino mass matrix as,

$$m_\nu^{\text{LISS}} = \begin{pmatrix} \frac{a^2\mu_{ee}}{M_{11}^2} & 0 & 0 \\ 0 & 0 & \frac{bc\mu_{\mu\tau}}{M_{22}M_{33}} \\ 0 & \frac{bc\mu_{\mu\tau}}{M_{22}M_{33}} & 0 \end{pmatrix} \tag{3.10}$$

which delivers light neutrino mass eigenstates with degenerate eigenvalues $\left\{\frac{a^2\mu_{ee}}{M_{11}^2}, -\frac{bc\mu_{\mu\tau}}{M_{22}M_{33}}, \frac{bc\mu_{\mu\tau}}{M_{22}M_{33}}\right\}$ similar to the previous situation given in eq. (2.2). Since the mass matrices M_D and M are diagonal in structure, the non-degenerate light neutrino masses consistent with observed values Δm^2_{sol} and Δm^2_{atm} can be achieved by suitable modification in the μ matrix. The modification in the matrix structure of μ matrix can be implemented with the inclusion of extra terms in the μ matrix which may be either of off-diagonal or diagonal in nature. Therefore, the extra singlet scalar χ with non-zero $\text{U}(1)_{L_\mu - L_\tau}$ charge which was originally introduced for spontaneous symmetry breaking of $\text{U}(1)_{L_\mu - L_\tau}$ symmetry can remove this degeneracy without affecting the usual left-right symmetry. We call this scenario as 'Extended LRSM with Inverse Seesaw (ELISS)'. The introduction of χ allows additional Yukawa-like terms in the Lagrangian and now the total Lagrangian for ELISS scenario becomes,

$$\mathcal{L}_{\text{ELISS}} = \mathcal{L}_{\text{LISS}} + \mathcal{L}_\chi \tag{3.11}$$

where $\mathcal{L}_\chi$ is the correction terms to the LISS lagrangian due to the introduction of new scalar χ.

$\boldsymbol{\mathcal{L}_\chi}$ responsible for off-diagonal correction to $\boldsymbol{\mu}$ matrix: considering the extra scalar χ with $\text{U}(1)_{L_\mu - L_\tau}$ charge 1, the modified Lagrangian with Yukawa-like terms can be written as,

$$\mathcal{L}_\chi \supset \mu_{e\mu} S^T_{e_L} S_{\mu_L} \chi^* + \mu_{e\tau} S^T_{e_L} S_{\tau_L} \chi + \mu_{e\mu} S^T_{\mu_L} S_{e_L} \chi^* + \mu_{e\tau} S^T_{\tau_L} S_{e_L} \chi \tag{3.12}$$

which modifies the structure of the light neutrino mass matrix now looking like,

$$m_\nu^{\text{ELISS}} = \begin{pmatrix} \frac{a^2\mu_{ee}}{M_{11}^2} & \frac{ab\mu_{e\mu}}{M_{11}M_{22}} & \frac{ac\mu_{e\tau}}{M_{11}M_{33}} \\ \frac{ab\mu_{e\mu}}{M_{11}M_{22}} & 0 & \frac{bc\mu_{\mu\tau}}{M_{22}M_{33}} \\ \frac{ac\mu_{e\tau}}{M_{11}M_{33}} & \frac{bc\mu_{\mu\tau}}{M_{22}M_{33}} & 0 \end{pmatrix} \tag{3.13}$$

Now, if we consider M_D and M as constant identity mass matrices i.e., $M_D = a\mathbb{I}_{3\times3}$ and $M = M_{11}\mathbb{I}_{3\times3}$, then $m_\nu^{\text{ELISS}} \sim \mu$. Since light neutrino mass matrix can be diagonalised by U_{PMNS} matrix [1],

$$|U_{\text{PMNS}}| \approx \begin{pmatrix} 0.814 & 0.554 & 0.147 \\ 0.329 & 0.572 & 0.717 \\ 0.432 & 0.555 & 0.742 \end{pmatrix} \tag{3.14}$$

we can diagonalise μ by $U_{\rm PMNS}$ and rewrite the mass matrix as (considering the couplings $\mu_{e\mu} = \mu_{e\tau}$ for sake of simplicity),

$$m_\nu'^{\rm ELISS} = \begin{pmatrix} \frac{a^2\mu_{ee}}{M_{11}^2} & \frac{a^2\mu_{e\mu}}{M_{11}^2} & \frac{a^2\mu_{e\mu}}{M_{11}^2} \\ \frac{a^2\mu_{e\mu}}{M_{11}^2} & 0 & \frac{a^2\mu_{\mu\tau}}{M_{11}^2} \\ \frac{a^2\mu_{e\mu}}{M_{11}^2} & \frac{a^2\mu_{\mu\tau}}{M_{11}^2} & 0 \end{pmatrix} \tag{3.15}$$

whose corresponding eigenvalues are $\{\frac{-a^2\mu_{\mu\tau}}{M_{11}^2}, m_\nu'^{{\rm ELISS}a} \pm m_\nu'^{{\rm ELISS}b}\}$ with

$$m_\nu'^{{\rm ELISS}a} = \frac{a^2}{2M_{11}^2}(\mu_{ee} + \mu_{\mu\tau}),$$
$$m_\nu'^{{\rm ELISS}b} = \frac{a^2}{2M_{11}^2}\sqrt{\mu_{ee}^2 + 8\mu_{e\mu}^2 - 2\mu_{ee}\mu_{\mu\tau} + \mu_{\mu\tau}^2}$$

$\mathcal{L}_\chi$ responsible for diagonal correction to μ matrix: similarly, if we consider χ with $\mathrm{U}(1)_{L_\mu - L_\tau}$ charge 2, then the Lagrangian can be written as,

$$\mathcal{L}_\chi \supset \mu_{\mu\mu} S_{\mu_L}^T S_{\mu_L} \chi^* + \mu_{\tau\tau} S_{\tau_L}^T S_{\tau_L} \chi \tag{3.16}$$

Now the modified light neutrino mass matrix in this framework can be expressed as,

$$m_\nu^{\rm ELISS} = \begin{pmatrix} \frac{a^2\mu_{ee}}{M_{11}^2} & 0 & 0 \\ 0 & \frac{b^2\mu_{\mu\mu}}{M_{22}^2} & \frac{bc\mu_{\mu\tau}}{M_{22}M_{33}} \\ 0 & \frac{bc\mu_{\mu\tau}}{M_{22}M_{33}} & \frac{c^2\mu_{\tau\tau}}{M_{33}^2} \end{pmatrix} \tag{3.17}$$

with mass eigenvalues $\{\frac{a^2\mu_{ee}}{M_{11}^2}, m_\nu^{\rm ELISS\ a} \pm m_\nu^{\rm ELISS\ b}\}$ where,

$$m_\nu^{\rm ELISS\ a} = \frac{c^2 M_{22}^2 \mu_{\tau\tau} + b^2 M_{33}^2 \mu_{\mu\mu}}{2M_{22}^2 M_{33}^2},$$
$$m_\nu^{\rm ELISS\ b} = \frac{\sqrt{c^4 M_{22}^4 \mu_{\tau\tau}^2 - 2b^2c^2 M_{22}^2 M_{33}^2 \mu_{\mu\mu}\mu_{\tau\tau} + b^4 M_{33}^4 \mu_{\mu\mu}^2 + 4b^2c^2 M_{22}^2 M_{33}^2 \mu_{\mu\tau}^2}}{2M_{22}^2 M_{33}^2}$$

We found that, both the cases i.e with the diagonal as well as off-diagonal corrections to μ-matrix successfully explain current-day neutrino oscillation data at 3σ interval of global fit by NuFIT 4.1 [93].

3.1 Non-standard neutrino interaction via non-unitarity effects in LISS

With the presence of extra sterile neutrinos on top of SM light active neutrinos, the measure of deviation of the neutrino mixing matrix from unitarity is known as non-unitarity effects in leptonic sector which can provide a new window to probe physics beyond Standard Model at present and planned neutrino factories. In the considered framework, left-right inverse seesaw scheme gives non-unitarity effects and thereby can generate non-standard neutrino interaction (NSI). The non-unitarity mixing matrix ($\mathbb{N}$) and the measure of deviation from unitarity (η) can be read as,

$$\begin{aligned}\mathbb{N} &\simeq \left(1-\frac{1}{2}\Theta\Theta^\dagger\right) U = (1-\eta)\, U\,,\\ \eta &= \frac{1}{2}\Theta\Theta^\dagger\,, \text{with} \quad \Theta \simeq M_D/M\end{aligned} \tag{3.18}$$

where, $U = U_{\rm PMNS}$ being the unitary matrix diagonalizing m_ν and $\Theta \simeq M_D/M$, in turn related to the light-heavy neutrino mixing matrix element $V^{\nu\xi}$ to be used next section onwards. This parameter $V^{\nu\xi}$ is a measure of how active light neutrinos mix with heavy sterile neutrinos and can be constrained from non-unitarity effects, NSI effects and muon anomalous $g-2$ etc. The present experimental bounds on the unitarity violation in $e\mu$, $e\tau$, $\mu\tau$, $\tau\tau$ sectors are $|\eta_{e\mu}| < 3.5\times10^{-5}$, $|\eta_{e\tau}| < 8.0\times10^{-4}$, $|\eta_{\mu\tau}| < 5.1\times10^{-3}$ and $|\eta_{\tau\tau}| < 2.7\times10^{-3}$ [98, 106] respectively. For a more detailed study on low energy LFV processes $\mu\to e\gamma$, $\mu\to eee$ and $\mu\to e$ conversion in nuclei due to non-unitarity effects readers can refer [107, 108]. Through charged current interaction and expressing light active neutrinos in terms of mass eigenstates including light as well as sterile neutrinos, the heavy sterile neutrinos (ξ) couple to gauge sector of SM which eventually create non-standard interaction (NSI) for neutrinos. In NSI effects for a given non-unitarity lepton mixing matrix $\mathbb{N}$, the vacuum neutrino oscillation probability $P_{\alpha\beta}$ can be expressed as [109],

$$\begin{aligned}P_{\alpha\beta} &= \sum_{i,j}\mathcal{F}^i_{\alpha\beta}\mathcal{F}^{j*}_{\alpha\beta} - 4\sum_{i>j}\mathrm{Re}(\mathcal{F}^i_{\alpha\beta}\mathcal{F}^{j*}_{\alpha\beta})\sin^2\left(\frac{\Delta m^2_{ij}L}{4E}\right)\\ &\quad +2\sum_{i>j}\mathrm{Im}(\mathcal{F}^i_{\alpha\beta}\mathcal{F}^{j*}_{\alpha\beta})\sin\left(\frac{\Delta m^2_{ij}L}{2E}\right)\,,\end{aligned} \tag{3.19}$$

where $\Delta m^2_{ij} \equiv m^2_i - m^2_j$ are the neutrino mass-squared differences and $\mathcal{F}^i$ are defined by

$$\mathcal{F}^i_{\alpha\beta} \equiv \sum_{\gamma,\rho}(R^*)_{\alpha\gamma}(R^*)^{-1}_{\rho\beta}U^*_{\gamma i}U_{\rho i}\,. \tag{3.20}$$

Here, the normalized non-unitary factor in terms of η parameters is given by

$$R_{\alpha\beta} \equiv \frac{(1-\eta)_{\alpha\beta}}{[(1-\eta)(1-\eta^\dagger)]_{\alpha\alpha}}\,. \tag{3.21}$$

The mass parameters M_D and M are found to be diagonal in the present framework with $\mathrm{U}(1)_{L_\mu-L_\tau}$ symmetry. This makes few elements of η negligible and hence restricts $R_{\alpha\beta}$. The neutrino factory will provide excellent sensitivity to probe these non-standard interaction effects and for more details about NSI in inverse seesaw mechanism, one may read refs [108–114]. Presence of heavy neutral fermions leads to non-unitarity effects and it has been shown that one may constrain these heavy fermions by studying their impact by adding few effective operators $\mathcal{O}^d$ of dimension $d > 4$ to the interaction Lagrangian [112]. We also skip the detailed phenomenology of LISS (interested readers may refer [96, 98–100, 109, 111–113, 115]) in the context of cLFV, non-unitarity effects, $0\nu\beta\beta$, LNV at collider. Rather, we intend to explore the implications of light-heavy neutrino mixing $V^{\nu\xi}$ with purely left-handed currents to new physics contributions to muon $g-2$ anomaly in the following section.

4 Prediction on muon $(g-2)$ anomaly

For a comprehensive review on new physics scenarios explaining muon $(g-2)$ anomaly one may refer [13, 14, 116]. Most of these works predict that new light gauge bosons and light neutral scalars are good candidates for addressing the anomaly since they contribute positively to Δa_μ. In our model, new contributions to muon $(g-2)$ anomaly arise from the interactions of;

- singly charged gauge bosons with heavy neutral fermions,
- neutral vector boson with singly charged fermions,
- singly charged scalars with neutral fermion,
- neutral scalars with muons,
- extra light new gauge boson $Z_{\mu\tau}$ with muons.

In the following we study analytically all these new physics contributions to Δa_μ and numerically estimate the individual contributions in the next section. Notably, for the calculation of Δa_μ we neglect the flavor mixing as they give negligible correction to the anomaly [116]. Another important point to recall here is that inverse seesaw mechanism which explains neutrino mass in this model also allows large light-heavy neutrino mixing due to which the contribution coming from the charged gauge boson interaction with heavy neutral fermion becomes sizeable.

4.1 Gauge boson contribution

Before moving on to the Feynman diagrams mediated by gauge bosons, we write the basic charged current(CC) interaction Lagrangian for leptons within left-right theories.

$$\mathcal{L}^l_{\rm cc} = \sum_{\alpha=e,\mu,\tau} \left[\frac{g_L}{\sqrt{2}}\bar{\ell}_{\alpha L}\gamma_\beta \ell_{\alpha L} W^\beta_L + \frac{g_R}{\sqrt{2}}\bar{\ell}_{\alpha R}\gamma_\beta \ell_{\alpha R} W^\beta_R\right] + \text{h.c.} \tag{4.1}$$

For Inverse Seesaw (ISS) mechanism [96, 98], the flavour eigenstates ν_L and N_R can be expressed in terms of admixture of mass eigenstates $(\nu_i,\, \xi_j)$ as follows,

$$\nu_{\mu L} = V^{\nu\nu}_{\mu i}\nu_i + V^{\nu\xi}_{\mu j}\xi_j \tag{4.2}$$

$$N_{\mu R} = V^{N\nu}_{\mu i}\nu_i + V^{N\xi}_{\mu j}\xi_j \tag{4.3}$$

where $i = 1, 2, 3$ goes over physical states for light neutrinos and $j = 1, 2, \ldots, 6$ runs over heavy states forming three pairs of pseudo-Dirac neutrinos. Using eq. (4.2) in the charged current interaction lagrangian given in eq. (4.1), we present the vector and axial vector couplings (g_v and g_a) in table 3.

Interaction Vertex	$g_{v2} = -g_{a2}$	Interaction Vertex	$g_{v1} = g_{a1}$
$\overline{\nu}_1 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 1}^{\nu\nu *}$	$\overline{\nu}_1 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 1}^{N\nu *}$
$\overline{\nu}_2 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 2}^{\nu\nu *}$	$\overline{\nu}_2 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 1}^{N\nu *}$
$\overline{\nu}_3 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 3}^{\nu\nu *}$	$\overline{\nu}_3 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 1}^{N\nu *}$
$\overline{\xi}_1 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 1}^{\nu\xi *}$	$\overline{\xi}_1 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 1}^{N\xi *}$
$\overline{\xi}_2 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 2}^{\nu\xi *}$	$\overline{\xi}_2 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 2}^{N\xi *}$
$\overline{\xi}_3 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 3}^{\nu\xi *}$	$\overline{\xi}_3 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 3}^{N\xi *}$
$\overline{\xi}_4 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 4}^{\nu\xi *}$	$\overline{\xi}_4 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 4}^{N\xi *}$
$\overline{\xi}_5 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 5}^{\nu\xi *}$	$\overline{\xi}_5 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 5}^{N\xi *}$
$\overline{\xi}_6 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 6}^{\nu\xi *}$	$\overline{\xi}_6 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 6}^{N\xi *}$

Table 3. Relevant vector and axial vector couplings for muon with W_L, W_R gauge bosons and physical neutral fermion states within the inverse seesaw (ISS) scenario.

In inverse seesaw scheme, the light neutrinos are Majorana in nature while heavy neutrinos are pseudo-Dirac. Alternatively, in extended inverse seesaw scenario (EISS) [98, 99] both light neutrino ν_L as well as heavy neutrinos S_L, N_R are purely Majorana in nature. Thus, the flavour eigenstates ν_L and N_R can be expressed in terms of admixture of mass eigenstates (ν_i, S_i, N_i) in the following way,

$$\nu_{\mu L} = V_{\mu i}^{\nu\nu} \nu_i + V_{\mu i}^{\nu S} S_i + V_{\mu i}^{\nu N} N_i \tag{4.4}$$

$$N_{\mu R} = V_{\mu i}^{N\nu} \nu_i + V_{\mu i}^{NS} S_i + V_{\mu i}^{NN} N_i \tag{4.5}$$

where $i = 1, 2, 3$ goes over physical states. For EISS we present the vector and axial vector couplings in table 4.

The diagrams in figure 1 are mediated by singly charged right-handed and left-handed gauge bosons W_R, W_L interacting with muons. Here ξ represents the heavy neutrino states in mass basis within the inverse seesaw framework. W_L can interact with heavy right-handed neutrino due to inverse seesaw mechanism in the model, and we find out in the next section that the most significant contribution to muon anomaly comes from this channel. For a detailed discussion on the contributions arising from singly charged vector bosons one may refer [117–120].

Figure 1(a): contribution due to W_R mediation; $\Delta a_\mu(\xi, W_R)$: for calculating its contribution, we start by sorting out the relevant interaction terms for this diagram.

$$\mathcal{L}_{\text{int}} = g_{v1} W_{R\mu}^+ \overline{\nu_\mu} \gamma^\mu \mu + g_{a1} W_{R\mu}^+ \overline{\nu_\mu} \gamma^\mu \gamma^5 \mu + \text{h.c.} \tag{4.6}$$

The contribution arising from this diagram to the anomalous magnetic moment can be determined by the following expression.

$$\Delta a_\mu(\xi, W_R) \simeq \frac{1}{8\pi^2} \frac{m_\mu^2}{m_{W_R}^2} \int_0^1 dx \frac{g_{v1}^2 P_{v1}(x) + g_{a1}^2 P_{a1}(x)}{\epsilon^2 \lambda^2 (1-x)(1-\epsilon^{-2} x) + x} \tag{4.7}$$

Interaction Vertex	$g_{v2} = -g_{a2}$	Interaction Vertex	$g_{v1} = g_{a1}$
$\overline{\nu}_1 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 1}^{\nu\nu *}$	$\overline{\nu}_1 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 1}^{N\nu *}$
$\overline{\nu}_2 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 2}^{\nu\nu *}$	$\overline{\nu}_2 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 1}^{N\nu *}$
$\overline{\nu}_3 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 3}^{\nu\nu *}$	$\overline{\nu}_3 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 1}^{N\nu *}$
$\overline{S}_1 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 1}^{\nu S *}$	$\overline{S}_1 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 1}^{N S *}$
$\overline{S}_2 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 2}^{\nu S *}$	$\overline{S}_2 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 2}^{N S *}$
$\overline{S}_3 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 3}^{\nu S *}$	$\overline{S}_3 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 3}^{N S *}$
$\overline{N}_1 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 1}^{\nu N *}$	$\overline{N}_1 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 1}^{N N *}$
$\overline{N}_2 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 2}^{\nu N *}$	$\overline{N}_2 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 2}^{N N *}$
$\overline{N}_3 \mu W_L^+$	$\frac{g_L}{2\sqrt{2}} V_{\mu 3}^{\nu N *}$	$\overline{N}_3 \mu W_R^+$	$\frac{g_R}{2\sqrt{2}} V_{\mu 3}^{N N *}$

Table 4. Relevant vector and axial vector couplings for muon with W_L, W_R gauge bosons and physical neutral fermion states within the extended inverse seesaw (EISS) scenario.

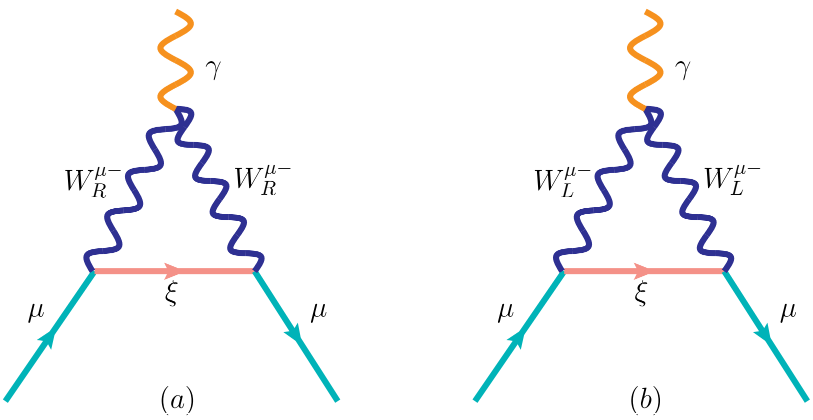

Figure 1. Feynman diagrams for the interaction of singly charged vector bosons: in left-panel due to the mediation of singly charged right-handed gauge boson W_R with heavy neutrinos and in right-panel due to the mediation of singly charged left-handed gauge boson W_L with exchange of heavy neutrinos. The W_L mediated diagram with exchange of heavy neutrinos gives sizeable contribution in ISS scheme.

where, m_μ is the mass of muon, m_{W_R} is the mass of right-handed charged gauge boson W_R, $\epsilon \equiv \left(\frac{m_{\nu_\mu}}{m_\mu}\right)$, $\lambda \equiv \left(\frac{m_\mu}{m_{W_R}}\right)$, and

$$P_{v1}(x) = 2x^2(1 + x - 2\epsilon) - \lambda^2(1-\epsilon)^2 x(1-x)(x+\epsilon)$$
$$P_{a1}(x) = 2x^2(1 + x + 2\epsilon) - \lambda^2(1+\epsilon)^2 x(1-x)(x-\epsilon)$$

After simplifying the integration the expression can be rewritten as (we will neglect the terms containing ϵ and λ in the expression of muon anomaly Δa_μ onwards (except neutral scalar sector to be discussed in next subsections) as they are really tiny corrections),

$$\Delta a_\mu(\xi, W_R) \simeq \frac{1}{4\pi^2}\frac{m_\mu^2}{m_{W_R}^2}\left[|g_{v1}^\mu|^2\left(\frac{5}{6}\right)+|g_{a1}^\mu|^2\left(\frac{5}{6}\right)\right]; \quad \text{with} \quad m_{W_R} \gg m_\mu. \tag{4.8}$$

Here we have, $|g_{v1}| = |g_{a1}| = \frac{g_R}{2\sqrt{2}}$ (as given in table 3 with $\mathcal{O}(1)$ neutrino mixing) and with these values we can rewrite eq. (4.8) as,

$$\Delta a_\mu(W_R) \simeq 2.3\times 10^{-11}\left(\frac{g_R}{g_L}\right)^2\left(\frac{1\ \text{TeV}}{m_{W_R}}\right)^2 \sum_{i=1,\ldots,6}|V_{\mu i}^{N\xi}|^2 \tag{4.9}$$

Figure 1(b): contribution due to W_L mediation with light-heavy neutrino mixing; $\Delta a_\mu(\xi, W_L)$: similar as 1(a) the relevant interaction terms for this diagram.

$$\mathcal{L}_{\text{int}} = g_{v1}W_{L\mu}^+\overline{\nu_\mu}\gamma^\mu\mu + g_{a1}W_{L\mu}^+\overline{\nu_\mu}\gamma^\mu\gamma^5\mu + \text{h.c.} \tag{4.10}$$

So, for W_L interacting with heavy neutrino the contribution to muon anomalous magnetic moment can be expressed as,

$$\Delta a_\mu(\xi, W_L) \simeq \frac{1}{4\pi^2}\frac{m_\mu^2}{m_{W_L}^2}\left[|g_{v2}^\mu|^2\left(\frac{5}{6}\right)+|g_{a2}^\mu|^2\left(\frac{5}{6}\right)\right]; \quad \text{with} \quad m_{W_L} \gg m_\mu. \tag{4.11}$$

Using the couplings for this interaction given in table 3 we can rewrite eq. (4.11) as,

$$\Delta a_\mu(\xi, W_L) \simeq 9.06\times 10^{-9}\ g_L^2 \sum_{i=1,\ldots,6}|V_{\mu i}^{\nu\xi}|^2 \tag{4.12}$$

Since the ISS scenario allows large mixing between light and heavy neutrinos, moving from flavor to mass basis we can see that for $\mathcal{O}(0.1)$ light-heavy neutrino mixing, heavy neutrinos with mass $\sim$ few GeV play a significant role in context of muon $g-2$ anomaly by interacting with W_L. Also, in the next section we will see that this gives positive and significant contribution to Δa_μ.

Figure 2: contribution due to Z_R mediation; $\Delta a_\mu(Z_R)$: the new contribution for muon anomalous $g-2$ arising from exchange of right-handed neutral gauge boson Z_R, as shown in figure 2, is derived from the neutral current interaction as,

$$\bar{\mu}\gamma_\beta\partial^\beta\mu + i\frac{g_L}{\sqrt{1-\delta\tan^2\theta_W}}\bar{\mu}\gamma_\beta(g_v - g_a\gamma^5)\mu Z_R^\beta \tag{4.13}$$

with the couplings

$$g_v = \frac{1}{4}\left[3\delta\tan^2\theta_W - 1\right]$$
$$g_a = \frac{1}{4}\left[1-\delta\tan^2\theta_W\right]$$

where $\delta = \frac{g_L^2}{g_R^2}$ and θ_W is the Weinberg angle. The Lagrangian for the charged fermions which interact with the SM leptons via a neutral vector boson (Z_R) can be written as

$$\mathcal{L}_{\text{int}} = g_{v3}Z_{R\mu}\overline{\mu}\gamma^\mu\mu + g_{a3}Z_{R\mu}\overline{\mu}\gamma^\mu\gamma^5\mu + \text{h.c.} \tag{4.14}$$

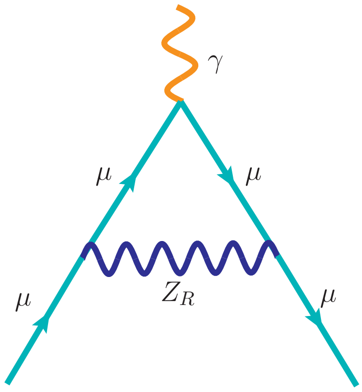

Figure 2. Feynman diagram for muon anomalous $g-2$ contribution arising from the mediation of right-handed neutral gauge boson Z_R with muons.

Using eq. (4.14) the contribution arising from Z_R to the muon anomalous magnetic moment can be expressed as,

$$\Delta a_\mu(Z_R) \simeq \frac{1}{8\pi^2}\frac{m_\mu^2}{m_{Z_R}^2}\int_0^1 dx\,\frac{g_{v3}^2 P_{v3}(x) + g_{a3}^2 P_{a3}(x)}{(1-x)(1-\lambda^2 x)+\lambda^2 x} \tag{4.15}$$

with $\lambda \equiv \left(\frac{m_\mu}{m_{Z_R}}\right)$, and

$$P_{v3}(x) = 2x^2(1-x)$$
$$P_{a3}(x) = 2x(1-x)(x-4) - 4\lambda^2 x^3$$

By simplifying the integrations the contribution is found to be,

$$\Delta a_\mu(Z_R) \simeq -\frac{1}{4\pi^2}\frac{m_\mu^2}{m_{Z_R}^2}\left[\left(-\frac{1}{3}\right)|g_{v3}^\mu|^2 + \left(\frac{5}{3}\right)|g_{a3}^\mu|^2\right]; \quad \text{with} \quad m_{Z_R} \gg m_\mu. \tag{4.16}$$

where the couplings g_{v3}, g_{a3} are same as g_v, g_a respectively as in eq. (4.13) and depending on the values of these vector and axial couplings the contribution can be either positive or negative.

4.2 Scalar sector contribution

The Yukawa Lagrangian involving scalars can be written as,

$$\mathcal{L}_{\text{Yuk}} = \bar{\ell}_L(Y_{22}\Phi + \tilde{Y}_{22}\tilde{\Phi})\ell_R + \bar{\ell}_R(Y_{22}\Phi^* + \tilde{Y}_{22}\tilde{\Phi}^*)\ell_L \tag{4.17}$$

where the scalar bidoublet Φ contains two charged scalars h_3^-, h_4^-, two neutral CP-even scalars h_1^0, h_2^0 and two neutral CP-odd scalars ϕ_1^0, ϕ_2^0 as

$$\Phi = \begin{pmatrix} v_1 + h_1^0 + i\phi_1^0 & h_3^+ \\ h_4^- & v_2 + h_2^0 + i\phi_2^0 \end{pmatrix}$$

and

$$\tilde{\Phi} = \sigma^2 \Phi^* \sigma^2 = \begin{pmatrix} v_2 + h_2^0 - i\phi_2^0 & -h_4^+ \\ -h_3^- & v_1 + h_1^0 - i\phi_1^0 \end{pmatrix}$$

The Feynman diagrams of these scalars interacting with muons are shown in figures 3, 4, 5 respectively. We later find out in section 5 that among these only the neutral CP-even scalars h_1^0, h_2^0 contribute positively to Δa_μ. Now by considering only muon family with

$$\ell_L = \begin{pmatrix} \nu_{\mu L} \\ \mu_L \end{pmatrix}, \quad \ell_R = \begin{pmatrix} N_{\mu R} \\ \mu_R \end{pmatrix},$$

the expanded Yukawa Lagrangian can be written as,

$$\begin{aligned} \mathcal{L}_{\rm Yuk} = & \left[\overline{\nu}_\mu \left[Y_{22}(v_1 + h_1^0 + i\phi_1^0) + \tilde{Y}_{22}(v_2 + h_2^0 - i\phi_2^0)\right] N_\mu + \overline{\nu}_\mu \left[Y_{22} h_3^+ - \tilde{Y}_{22} h_4^+\right] \mu\right] \frac{(1+\gamma_5)}{2} \\ & + \left[\overline{\mu} \left[Y_{22} h_4^- - \tilde{Y}_{22} h_3^-\right] N_\mu + \overline{\mu} \left[Y_{22}(v_2 + h_2^0 + i\phi_2^0) + \tilde{Y}_{22}(v_1 + h_1^0 - i\phi_1^0)\right] \mu\right] \frac{(1+\gamma_5)}{2} \\ & + \left[\overline{N}_\mu \left[Y_{22}(v_1 + h_1^0 - i\phi_1^0) + \tilde{Y}_{22}(v_2 + h_2^0 + i\phi_2^0)\right] \nu_\mu + \overline{N}_\mu \left[Y_{22} h_3^- - \tilde{Y}_{22} h_4^-\right] \mu\right] \frac{(1-\gamma_5)}{2} \\ & + \left[\overline{\mu} \left[Y_{22} h_4^+ - \tilde{Y}_{22} h_3^+\right] \nu_\mu + \overline{\mu} \left[Y_{22}(v_2 + h_2^0 - i\phi_2^0) + \tilde{Y}_{22}(v_1 + h_1^0 + i\phi_1^0)\right] \mu\right] \frac{(1-\gamma_5)}{2} \end{aligned} \tag{4.18}$$

The relevant terms in the Yukawa Lagrangian for the Feynman diagrams given in figure 3 are as follows,

$$\mathcal{L}_{\rm Yuk}(h_3^+, h_4^+) = \overline{\nu}_\mu \left[Y_{22} h_3^+ - \tilde{Y}_{22} h_4^+\right] \mu \frac{(1+\gamma_5)}{2} + \overline{\mu} \left[Y_{22} h_4^+ - \tilde{Y}_{22} h_3^+\right] \nu_\mu \frac{(1-\gamma_5)}{2} \tag{4.19}$$

The same equation can be written in mass basis using 4.2 as,

$$\begin{aligned} \mathcal{L}^{\rm mass}_{\rm Yuk}(h_3^+, h_4^+) = & [V_{\mu 1}^{\nu\nu *}\overline{\nu}_1 + V_{\mu 2}^{\nu\nu *}\overline{\nu}_2 + V_{\mu 3}^{\nu\nu *}\overline{\nu}_3 + V_{\mu 1}^{\nu S *}\overline{S}_1 + V_{\mu 2}^{\nu S *}\overline{S}_2 + V_{\mu 3}^{\nu S *}\overline{S}_3 + V_{\mu 1}^{\nu N *}\overline{N}_1 \\ & + V_{\mu 2}^{\nu N *}\overline{N}_2 + V_{\mu 3}^{\nu N *}\overline{N}_3] \left[Y_{22} h_3^+ - \tilde{Y}_{22} h_4^+\right] \mu \frac{(1+\gamma_5)}{2} + \overline{\mu} \left[Y_{22} h_4^+ - \tilde{Y}_{22} h_3^+\right] \\ & [V_{\mu 1}^{\nu\nu}\nu_1 + V_{\mu 2}^{\nu\nu}\nu_2 + V_{\mu 3}^{\nu\nu}\nu_3 + V_{\mu 1}^{\nu S} S_1 + V_{\mu 2}^{\nu S} S_2 + V_{\mu 3}^{\nu S} S_3 + V_{\mu 1}^{\nu N} N_1 + V_{\mu 2}^{\nu N} N_2 \\ & + V_{\mu 3}^{\nu N} N_3] \frac{(1-\gamma_5)}{2} \end{aligned} \tag{4.20}$$

The diagrams in figure 3 represent the interactions mediated by singly charged scalars h_3^- and h_4^-.

Figure 3(a): contribution due to charged scalar, h_3^+ mediation; $\Delta a_\mu(h_3^+)$: the relevant interaction terms involving singly charged scalar with scalar coupling (g_{s1}) and pseudo-scalar coupling (g_{p1}) are given by,

$$\mathcal{L}_{\rm int} = g_{s1} h_3^+ \overline{\nu_\mu} \mu + g_{p1} h_3^+ \overline{\nu_\mu} \gamma^5 \mu + \text{h.c.} \tag{4.21}$$

In general, the contribution of a singly charged scalar to the muon anomaly can be expressed as,

$$\Delta a_\mu(h_3^+) \simeq \frac{1}{8\pi^2} \frac{m_\mu^2}{m_{h_3^+}^2} \int_0^1 dx \frac{g_{s1}^2 P_{s1}(x) + g_{p1}^2 P_{p1}(x)}{\epsilon^2 \lambda^2 (1-x)(1-\epsilon^{-2} x) + x} \tag{4.22}$$

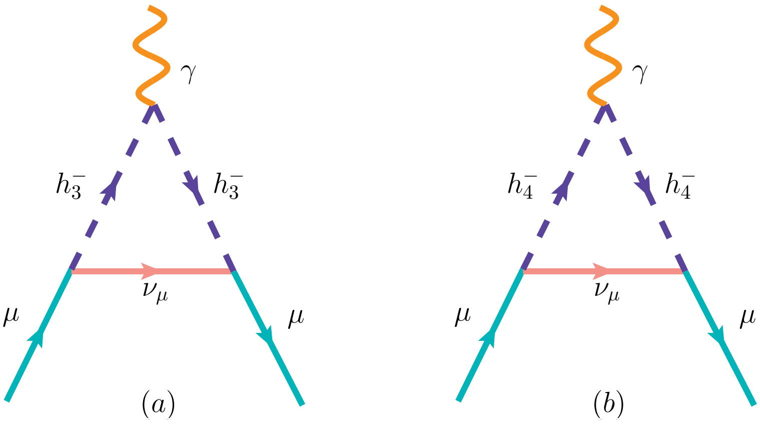

Figure 3. Feynman diagrams for the interaction of singly charged scalars h_3^-, h_4^- with muons contributing to the muon anomalous $g-2$.

with $\epsilon \equiv \left(\frac{m_{\nu_\mu}}{m_\mu}\right)$, $\lambda \equiv \left(\frac{m_\mu}{m_{h_3^+}}\right)$ and

$$P_{s1}(x) = -x(1-x)(x+\epsilon)$$
$$P_{p1}(x) = -x(1-x)(x-\epsilon)$$

So, in this case the extra contribution is found to be,

$$\Delta a_\mu(h_3^+) \simeq -\frac{1}{4\pi^2}\frac{m_\mu^2}{m_{h_3^+}^2}\left[|g_{s1}^\mu|^2\left(\frac{1}{12}\right) + |g_{p1}^\mu|^2\left(\frac{1}{12}\right)\right]; \quad \text{with} \quad m_{h_3^+} \gg m_\mu, m_{\nu_\mu} \tag{4.23}$$

Figure 3(b): contribution due to charged scalar, h_4^+ mediation; $\Delta a_\mu(h_4^+)$: similarly the interaction terms involving h_4^+ with scalar coupling (g_{s2}) and pseudo-scalar coupling (g_{p2}) are,

$$\mathcal{L}_{\rm int} = g_{s2}h_4^+\overline{\nu_\mu}\mu + g_{p2}h_4^+\overline{\nu_\mu}\gamma^5\mu + \text{h.c.} \tag{4.24}$$

The expression for the contribution arising from this scalar to the muon anomaly can be written as,

$$\Delta a_\mu(h_4^+) \simeq -\frac{1}{4\pi^2}\frac{m_\mu^2}{m_{h_4^+}^2}\left[|g_{s2}^\mu|^2\left(\frac{1}{12}\right) + |g_{p2}^\mu|^2\left(\frac{1}{12}\right)\right]; \quad \text{with} \quad m_{h_4^+} \gg m_\mu, m_{\nu_\mu} \tag{4.25}$$

The couplings for the above two cases can be found from eq. (4.20) and are given in table 5.

The diagrams in figure 4 are mediated by CP-even neutral scalars h_1^0 and h_2^0.

Figure 4(a): contribution due to CP-even scalar, h_1^0 mediation; $\Delta a_\mu(h_1^0)$: in general if extra electrically neutral scalar fields are present in a model, they induce a shift in the leptonic magnetic moments via the following interactions:

$$\mathcal{L}_{\rm int} = g_{s3}h_1^0\overline{\mu}\mu + ig_{p3}h_1^0\overline{\mu}\gamma^5\mu \tag{4.26}$$

Interaction Vertex	$g_{s1} = g_{p1}$	Interaction Vertex	$g_{s2} = g_{p2}$
$\overline{\nu}_1 \mu h_3^+$	$\frac{Y_{22}}{2} V_{\mu 1}^{\nu\nu *}$	$\overline{\nu}_1 \mu h_4^+$	$-\frac{\tilde{Y}_{22}}{2} V_{\mu 1}^{\nu\nu *}$
$\overline{\nu}_2 \mu h_3^+$	$\frac{Y_{22}}{2} V_{\mu 2}^{\nu\nu *}$	$\overline{\nu}_2 \mu h_4^+$	$-\frac{\tilde{Y}_{22}}{2} V_{\mu 2}^{\nu\nu *}$
$\overline{\nu}_3 \mu h_3^+$	$\frac{Y_{22}}{2} V_{\mu 3}^{\nu\nu *}$	$\overline{\nu}_3 \mu h_4^+$	$-\frac{\tilde{Y}_{22}}{2} V_{\mu 3}^{\nu\nu *}$
$\overline{S}_1 \mu h_3^+$	$\frac{Y_{22}}{2} V_{\mu 1}^{\nu S *}$	$\overline{S}_1 \mu h_4^+$	$-\frac{\tilde{Y}_{22}}{2} V_{\mu 1}^{\nu S *}$
$\overline{S}_2 \mu h_3^+$	$\frac{Y_{22}}{2} V_{\mu 2}^{\nu S *}$	$\overline{S}_2 \mu h_4^+$	$-\frac{\tilde{Y}_{22}}{2} V_{\mu 2}^{\nu S *}$
$\overline{S}_3 \mu h_3^+$	$\frac{Y_{22}}{2} V_{\mu 3}^{\nu S *}$	$\overline{S}_3 \mu h_4^+$	$-\frac{\tilde{Y}_{22}}{2} V_{\mu 3}^{\nu S *}$
$\overline{N}_1 \mu h_3^+$	$\frac{Y_{22}}{2} V_{\mu 1}^{\nu N *}$	$\overline{N}_1 \mu h_4^+$	$-\frac{\tilde{Y}_{22}}{2} V_{\mu 1}^{\nu N *}$
$\overline{N}_2 \mu h_3^+$	$\frac{Y_{22}}{2} V_{\mu 2}^{\nu N *}$	$\overline{N}_2 \mu h_4^+$	$-\frac{\tilde{Y}_{22}}{2} V_{\mu 2}^{\nu N *}$
$\overline{N}_3 \mu h_3^+$	$\frac{Y_{22}}{2} V_{\mu 3}^{\nu N *}$	$\overline{N}_3 \mu h_4^+$	$-\frac{\tilde{Y}_{22}}{2} V_{\mu 3}^{\nu N *}$

Table 5. Relevant couplings associated with the Feynman diagrams involving h_3^-, h_4^- given in figure 3.

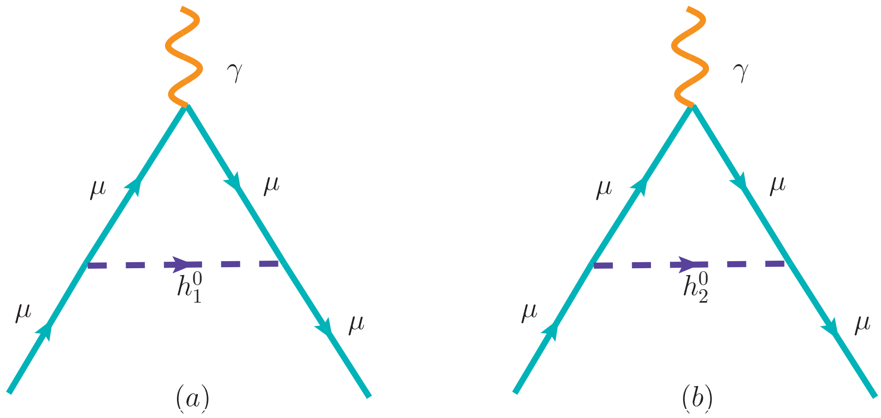

Figure 4. Feynman diagrams for the interaction of neutral CP-even scalars h_1^0, h_2^0 with muons.

From eq. (4.26) one can see that scalar and pseudo-scalar couplings shift $(g-2)_\mu$ by

$$\Delta a_\mu(h_1^0) \simeq \frac{1}{4\pi^2} \frac{m_\mu^2}{m_{h_1^0}^2} \int_0^1 dx \frac{g_{s3}^2 P_{s3}(x) + g_{p3}^2 P_{p3}(x)}{(1-x)(1-\lambda^2 x) + \lambda^2 x} \tag{4.27}$$

with $\lambda \equiv \left(\frac{m_\mu}{m_{h_1^0}}\right)$ and $P_{s3}(x) = x^2(2-x)$, $P_{p3}(x) = -x^3$.

So, from here we have the extra contribution to the anomalous magnetic moment as,

$$\Delta a_\mu(h_1^0) \simeq \frac{1}{4\pi^2} \frac{m_\mu^2}{m_{h_1^0}^2} \left[|g_{s3}^\mu|^2 \left(-\frac{7}{12} - \log\lambda \right) + |g_{p3}^\mu|^2 \left(\frac{11}{12} + \log\lambda \right) \right]; \quad \text{with} \quad m_{h_1^0} \gg m_\mu. \tag{4.28}$$

JHEP09(2020)010

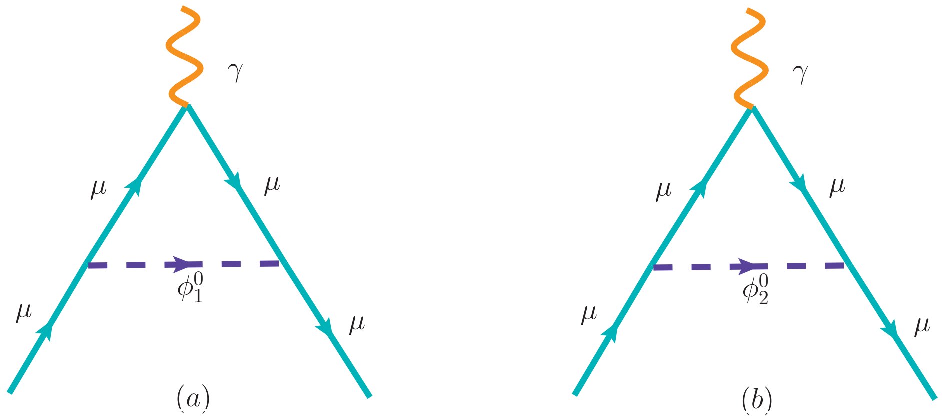

Figure 5. Feynman diagrams for the interaction of neutral CP-odd scalars ϕ_1^0, ϕ_2^0 with muons.

The result in eq. (4.28) is for general neutral scalars with scalar and pseudo-scalar couplings in the regime $m_{\text{Neutral Scalar}} \gg m_\mu$. The contribution coming from pure scalar can be derived from eq. (4.28) by setting the pseudo-scalar coupling (g_p) to zero and that from pseudo-scalar by setting the scalar coupling (g_s) to zero. By comparing with eq. (4.18) we have the couplings $g_{s3} = \tilde{Y}_{22},\ g_{p3} = 0$.

Figure 4(b): contribution due to CP-even scalar, h_2^0 mediation; $\Delta a_\mu(h_2^0)$: for this diagram the interaction Lagrangian can be written as

$$\mathcal{L}_{\text{int}} = g_{s4} h_2^0 \overline{\mu}\mu + i g_{p4} h_2^0 \overline{\mu}\gamma^5 \mu \tag{4.29}$$

Similar to the previous case its contribution to the anomalous magnetic moment can be written as,

$$\Delta a_\mu(h_2^0) \simeq \frac{1}{4\pi^2}\frac{m_\mu^2}{m_{h_2^0}^2}\left[|g_{s4}^\mu|^2\left(-\frac{7}{12}-\log\lambda\right)+|g_{p4}^\mu|^2\left(\frac{11}{12}+\log\lambda\right)\right]; \quad \text{with} \quad m_{h_2^0} \gg m_\mu. \tag{4.30}$$

From comparison with eq. (4.18) the couplings are $g_{s4} = Y_{22},\ g_{p4} = 0$.

Figure 5(a): contribution due to CP-odd scalar, ϕ_1^0 mediation; $\Delta a_\mu(\phi_1^0)$: in this case the interaction Lagrangian is given by,

$$\mathcal{L}_{\text{int}} = g_{s5} \phi_1^0 \overline{\mu}\mu + i g_{p5} \phi_1^0 \overline{\mu}\gamma^5 \mu \tag{4.31}$$

As in the case 4(a), here we will have the extra contribution to the anomalous magnetic moment as,

$$\Delta a_\mu(\phi_1^0) \simeq \frac{1}{4\pi^2}\frac{m_\mu^2}{m_{\phi_1^0}^2}\left[|g_{s5}^\mu|^2\left(-\frac{7}{12}-\log\lambda\right)+|g_{p5}^\mu|^2\left(\frac{11}{12}+\log\lambda\right)\right]; \quad \text{with} \quad m_{\phi_1^0} \gg m_\mu \tag{4.32}$$

The couplings here are $g_{s5} = 0,\ g_{p5} = -\tilde{Y}_{22}$.

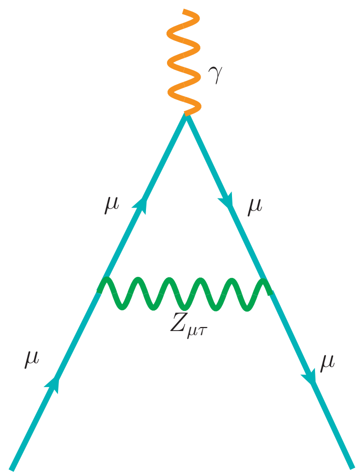

Figure 6. Feynman diagram for the interaction of new light gauge boson $Z_{\mu\tau}$ with muons.

Figure 5(b): contribution due to CP-odd scalar, $\boldsymbol{\phi_2^0}$ mediation; $\boldsymbol{\Delta a_\mu(\phi_2^0)}$: for this interaction the Lagrangian can be written as,

$$\mathcal{L}_{\text{int}} = g_{s6}\phi_2^0\overline{\mu}\mu + ig_{p6}\phi_2^0\overline{\mu}\gamma^5\mu \tag{4.33}$$

and its contribution to Δa_μ is,

$$\Delta a_\mu(\phi_2^0) \simeq \frac{1}{4\pi^2}\frac{m_\mu^2}{m_{\phi_2^0}^2}\left[|g_{s6}^\mu|^2\left(-\frac{7}{12}-\log\lambda\right)+|g_{p6}^\mu|^2\left(\frac{11}{12}+\log\lambda\right)\right]; \quad \text{with} \quad m_{\phi_2^0} \gg m_\mu \tag{4.34}$$

The couplings for this case are $g_{s6} = 0, \; g_{p6} = Y_{22}$.

Figure 6: contribution due to extra neutral gauge boson, $\boldsymbol{Z_{\mu\tau}}$ mediation; $\boldsymbol{\Delta a_\mu(Z_{\mu\tau})}$: this diagram 6 comes from the interaction of the new gauge boson $Z_{\mu\tau}$ associated with $\mathrm{U}(1)_{L_\mu - L_\tau}$ symmetry with muons. We have the terms in the Lagrangian

$$\sum_{\alpha=e,\mu,\tau}\left[\overline{\ell}_{\alpha L}\gamma^\mu D_\mu \ell_{\alpha L} + \overline{\ell}_{\alpha R}\gamma^\mu D_\mu \ell_{\alpha R}\right] \tag{4.35}$$

with covariant derivative $D_\beta = \partial_\beta + ig_{\mu\tau}qZ_\beta^{\mu\tau}$, where $g_{\mu\tau}$ is the gauge coupling of $\mathrm{U}(1)_{L_\mu - L_\tau}$ symmetry and q is the corresponding $L_\mu - L_\tau$ charge ($q_{\mu,\nu_\mu} = 1, q_{\tau,\nu_\tau} = -1$). By expanding this term explicitly for μ-family we will get $g_{\mu\tau}\overline{\mu}\gamma^\beta\mu Z_\beta^{\mu\tau}$ and this term contributes to muon $(g-2)$ anomaly.
So, the interaction Lagrangian can be written as,

$$\mathcal{L}_{\text{int}} = g_{\mu\tau}Z_{\mu\tau}\overline{\mu}\gamma^\mu\mu \tag{4.36}$$

Defining the parameter $\lambda \equiv \left(\frac{m_\mu}{m_{Z_{\mu\tau}}}\right)$, its contribution to the anomaly can be written as,

$$\Delta a_\mu(Z_{\mu\tau}) \simeq \frac{g_{\mu\tau}^2}{8\pi^2}\frac{m_\mu^2}{m_{Z_{\mu\tau}}^2}\int_0^1 dx \frac{2x^2(1-x)}{(1-x)(1-\lambda^2 x)+\lambda^2 x} \tag{4.37}$$

After simplifying the integrations its contribution can be written as,

$$\Delta a_\mu(Z_{\mu\tau}) = \frac{g_{\mu\tau}^2}{12\pi^2}\frac{m_\mu^2}{m_{Z_{\mu\tau}}^2}; \quad \text{with} \quad \lambda \equiv \left(\frac{m_\mu}{m_{Z_{\mu\tau}}}\right). \tag{4.38}$$

5 Results and discussion

Using the analytical expressions for different Feynman diagrams given in section 4, we plot the dependence of Δa_μ on the masses of the various species. For the purpose of understanding the behaviour we retain a large range for each of the mass values, although as we see much of it is excluded by the collider data. The excluded regions are clearly marked out. We see that the contribution of each of the class of diagrams independently could explain the entire anomaly, however for several of the species the mass value that would have allowed this is already ruled out.

We use the data $m_{W_L} = 80.4\,\text{GeV}, m_Z = 91.2\,\text{GeV}$, while for charge scalars we use the bounds [1],

$$m_{h^+} > 181\,\text{GeV}, m_{h^0,\phi^0} > 389\,\text{GeV}. \tag{5.1}$$

For the standard results in the graphs the dashed green line represents the current bound on Δa_μ. The red dashed lines represent the current 1σ bound on Δa_μ. The values of these standard results [116] are given below.

$$\Delta a_\mu(\text{Current Bound}) = (295 \pm 81) \times 10^{-11}$$
$$\Delta a_\mu(1\sigma \text{ Current Bound}) = 81 \times 10^{-11}$$

It is useful to keep in mind the projected and 1σ bounds on Δa_μ contribution which are $(295 \pm 34) \times 10^{-11}$ and 34×10^{-11} respectively, since they may be soon reached, though we have not used them in our plots.

We include in our analysis the important possibility of asymmetric LRSM. Usually in a left-right symmetric theory the $\mathrm{SU}(2)_L$ and $\mathrm{SU}(2)_R$ gauge couplings are equal, i.e. $g_L = g_R$, known as *symmetric LRSM* scenario. But, there is also the possibility that the Parity symmetry breaks at a higher scale than the $\mathrm{SU}(2)_R$ gauge symmetry, in which case the left-handed and right-handed gauge couplings become unequal, i.e. $g_L \neq g_R$. Such a model is called *asymmetric LRSM*, which was first proposed in [121] and more about this can be found in [122–127]. Hence, we have considered two different cases based on g_L and g_R for calculating the contributions of right-handed vector bosons W_R and Z_R to Δa_μ.

$$\text{Case I:} \qquad g_L = g_R = 0.653$$
$$\text{Case II:} \qquad g_L = 0.653, \quad g_R = 0.39$$

where the latter case corresponds to Pati-Salam breaking scale of 10^6 GeV, with grand unification in SO(10) at $10^{17.2}$ GeV [127].

With these representative values for gauge couplings we have numerically estimated and tabulated the upper bound on the muon anomaly contributions due to W_R and Z_R

Particles	Bounds on masses of mediators	$g_L = g_R$ (**Case I**)	$g_L \neq g_R$ (**Case II**)
$\Delta a_\mu(W_R)$	$\geq 4.1\,\text{TeV}$ [128]	$\leq 1.45 \times 10^{-12}$	$\leq 0.55 \times 10^{-12}$
$\Delta a_\mu(Z_R)$	$\geq 4.9\,\text{TeV}$ $(g_L = g_R)$ & $\geq 9.0\,\text{TeV}$ $(g_L \neq g_R)$	$\leq -0.64 \times 10^{-12}$	$\leq 0.10 \times 10^{-12}$

Table 6. Estimated values of the individual contributions coming from W_R and Z_R in extended LRSM for the cases $g_L = g_R$ and $g_L \neq g_R$. Relation between M_{W_R} and M_{Z_R} in LRSM with Higgs doublets can be found in ref. [87], from which we have obtained the lower bound on M_{Z_R} for our framework.

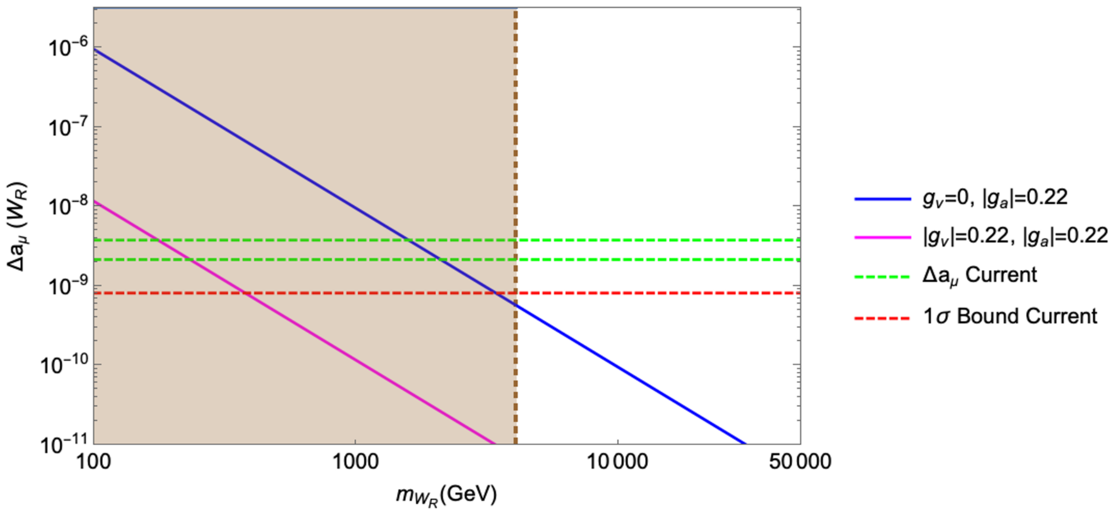

Figure 7. Plot showing the contribution of charged vector boson W_R to Δa_μ for the case $g_L = g_R$. The blue line represents the contribution of W_R when purely axial vector-like coupling is considered and magenta line represents the contribution when both vector-like and axial vector-like couplings are considered non-zero. It shows with purely axial vector-like coupling W_R with mass 2 TeV can address the anomaly whereas the case with combination of both couplings can satisfy the muon anomaly bound for $m_{W_R} \sim 200\,\text{GeV}$. Brown shaded region indicates the excluded mass range of W_R from collider constraints. So both the muon anomaly contributions is ruled out from collider constraints on m_{W_R}.

mediation channels (from their lower mass bound) for symmetric as well as asymmetric LRSM scenario in table 6.

The contributions arising from charged gauge boson W_R for the cases (i) $g_L = g_R$ (symmetric case), (ii) $g_L \neq g_R$ (asymmetric case) are presented in figures 7 and 8 respectively.

(i) For the case $g_L = g_R$, if we consider purely axial-vector like coupling i.e. $|g_v| = 0$ and $|g_a| = 0.22$ then the gauge boson W_R with mass around 2 TeV can address the whole anomaly. This is represented by the blue solid line in figure 7. whereas if we consider non-zero values for both couplings; $|g_v| = 0.22$ and $|g_a| = 0.22$, then the mass of W_R lies around 200 GeV (magenta line). Thus both the cases fall in the excluded mass range of W_R (brown shaded region) from collider bound.

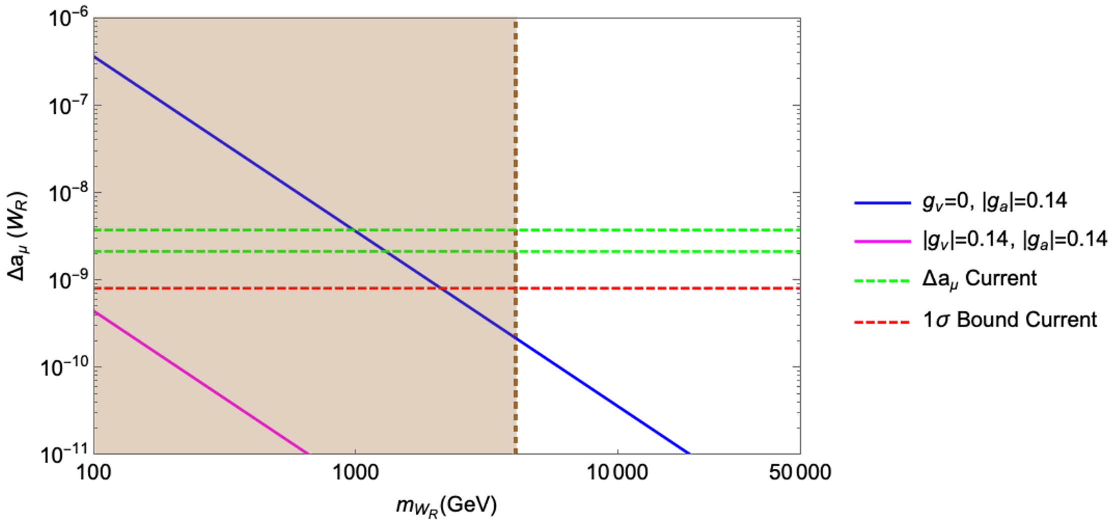

Figure 8. Plot showing the contribution of charged vector boson W_R to Δa_μ for the case $g_L \neq g_R$. The blue line represents the contribution of W_R when purely axial vector-like coupling is considered and it is sensitive to current bound on Δa_μ when W_R mass lies around 1 TeV. The magenta line represents the contribution when both vector-like and axial vector-like couplings are considered non-zero which is sensitive to the current bound on Δa_μ for mass range < 100 GeV but these two cases do not even satisfy the bound on W_R mass (as they both appear within the brown shaded excluded mass range).

(ii) Similarly for the case $g_L \neq g_R$, when purely axial-vector like coupling is considered i.e. $|g_v| = 0$ and $|g_a| = 0.14$ then W_R with mass around 1 TeV can explain the entire anomaly and the same is represented by blue line in figure 8. But for $|g_v| = 0.14$ and $|g_a| = 0.14$ the mass of W_R lies below 100 GeV (magenta line). This implies that even though W_R can explain the entire anomaly in both symmetric as well as asymmetric case, it is irrelevant for calculation since such a low mass for W_R is ruled out by collider experiments (brown shaded region) [128]. Though from the experimental side, where W_R interacts only with right handed neutrinos, i.e for $g_a = g_v$ the LEP bound on $\frac{g_v}{m_{W_R}}$ reads as $\frac{g_v}{m_{W_R}} < 4.8 \times 10^{-3}\,\text{GeV}^{-1}$ [129]. In our case $\frac{g_v}{m_{W_R}} \sim 5.1 \times 10^{-5}\,\text{GeV}^{-1}$ which clearly satisfies the bound.

Figures 9 and 10 show the contributions coming from the right-handed neutral vector boson Z_R for the cases $g_L = g_R$ and $g_L \neq g_R$ respectively. For the case $g_L = g_R$, Z_R gives negative contribution to Δa_μ and thus it is not relevant for our calculation, but for comparison perspective we have plotted the absolute value of the contributions vs Z_R mass in Log-Log plots. We have shown the excluded mass range for Z_R due to collider constraints as the brown shaded region in both these plots.

(i) For the case $g_L = g_R$, Z_R contributes positively and could have addressed the anomaly with $M_{Z_R} \sim 10$ GeV when purely vector-like contribution is considered (black line in figure 9), but which lies deep in the excluded region. The other two contributions

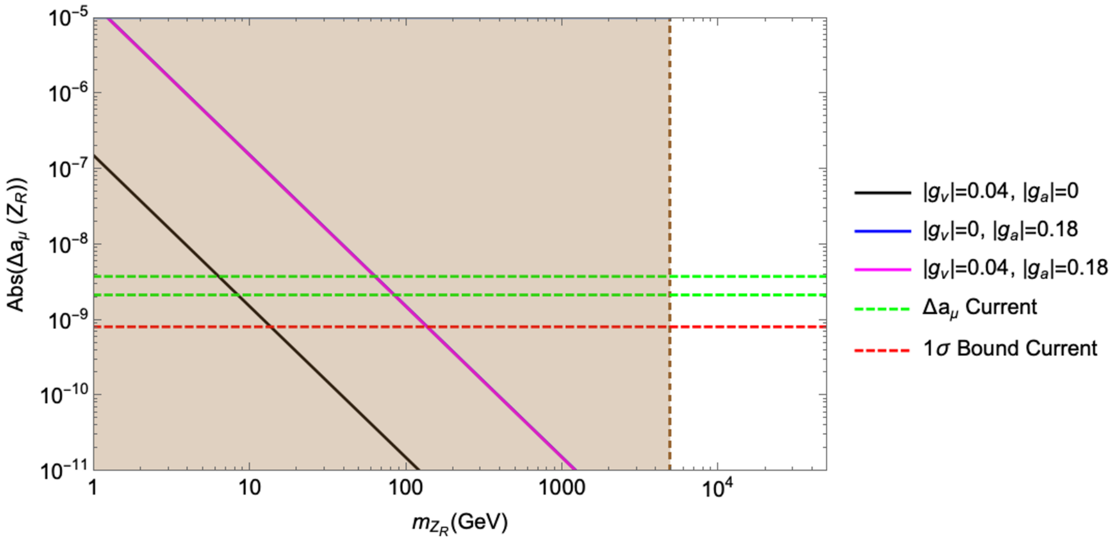

Figure 9. Plot showing the contribution of neutral vector boson Z_R to Δa_μ for the case $g_L = g_R$. The black line represents the contribution of Z_R when purely vector-like coupling is considered and the magenta line represents the contribution when both vector-like and axial vector-like couplings are considered non-zero. The contribution with purely axial vector-like coupling is negative so we have plotted the absolute value of it here and it merges with magenta line. Even though the contribution from Z_R is positive with purely vector-like coupling, all the cases fail to satisfy the current bound on Z_R mass (as recent collider developments exclude the brown shaded mass range for Z_R).

i.e. purely axial-vector like (blue line) and combination of both couplings (magenta line) give negative contributions.

(ii) For the case $g_L \neq g_R$, Z_R contributes positively for all the choices on couplings i.e. purely vector-like (black line in figure 10), purely axial-vector-like (blue line) as well as combination of both (magenta line) and can explain the anomaly but with $M_{Z_R} \sim 20-50\,\text{GeV}$, far below the collider bounds. This accords with ref. [129] which argues that a 95% C.L upper bound from LEP measurements applies for $g_v = g_a$ and $m_{Z_R} > \sqrt{s}$ that puts $\frac{g_v}{m_{Z_R}} < 2.2 \times 10^{-4}\,\text{GeV}^{-1}$ and thus discards the idea of a single Z_R boson explaining the anomaly. Some more bounds are given in refs. [129, 130].

Figure 11 shows the contributions coming from the charged scalars h_3^+, h_4^+ for three different choices of the couplings; $|g_s| = 0.29$ and $|g_p| = 0$ (purely scalar), $|g_s| = 0$ and $|g_p| = 0.29$ (purely pseudo-scalar) and $|g_s| = 0.29$ and $|g_p| = 0.29$ (combination of both). We have already discussed in section 4 that h_3^+, h_4^+ contribute negatively to Δa_μ, and thus we have plotted the absolute values of these contributions in Log-Log plot. Here the black and blue lines representing purely scalar and purely pseudo-scalar couplings which coincide together. Magenta line represents the contribution coming from the charged scalar sector when we consider both the scalar as well as pseudo-scalar couplings non-zero. The plot shows that the masses of the charged scalars lie around $\mathcal{O}(50)$ GeV which cannot satisfy the collider bounds on masses given in relation 5.1 (brown shaded region of the plot shows

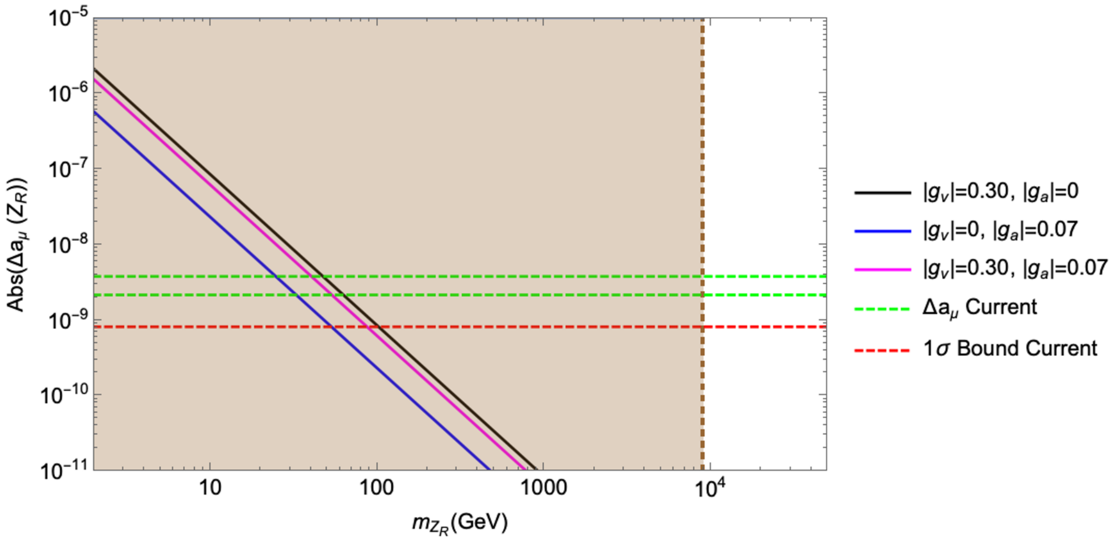

Figure 10. Plot showing the contribution of neutral vector boson Z_R to Δa_μ for the case $g_L \neq g_R$. The black line, blue line and magenta line represent the contributions of Z_R when purely vector-like, purely axial vector-like and the combination of both couplings are considered non-zero respectively. Even though all contributions are positive and sensitive to the current bounds on Δa_μ, none of the cases satisfy the bound on Z_R mass since all the mass values which can satisfy the muon anomaly bounds are residing in the brown shaded excluded mass range.

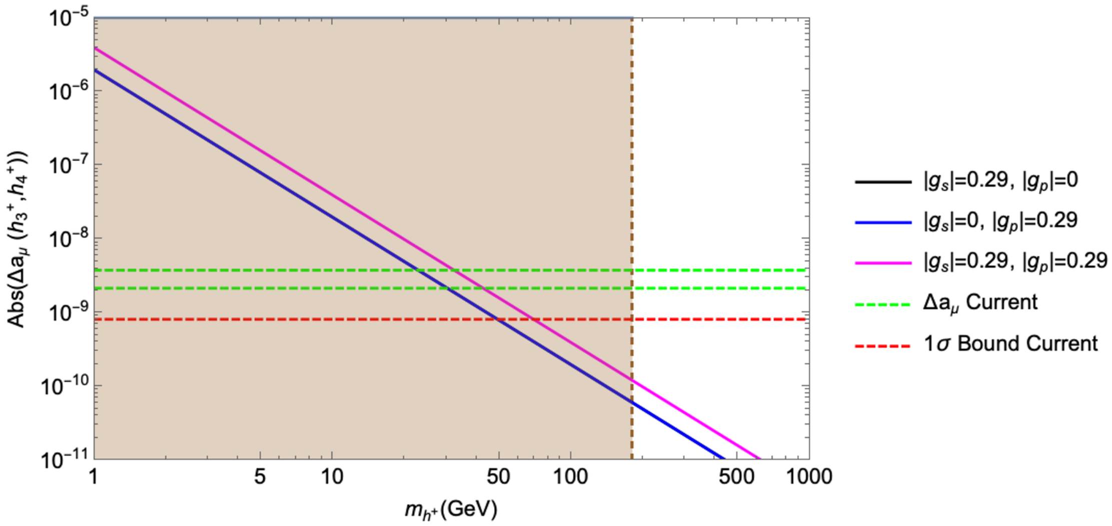

Figure 11. Plot showing the contributions of singly charged scalars h_3^+, h_4^+ to Δa_μ for three different choices of couplings; purely scalar, purely pseudo-scalar and combination of both. All the contributions are negative and thus are plotted their absolute values in log-log plot. The contributions coming from purely scalar and purely pseudo-scalar couplings are super-imposed and represented by the blue line. The magenta line represents the contribution from the combination of both couplings. Also the mass range for charged scalars $< 181\,\text{GeV}$ is ruled out from collider studies (brown shaded region).

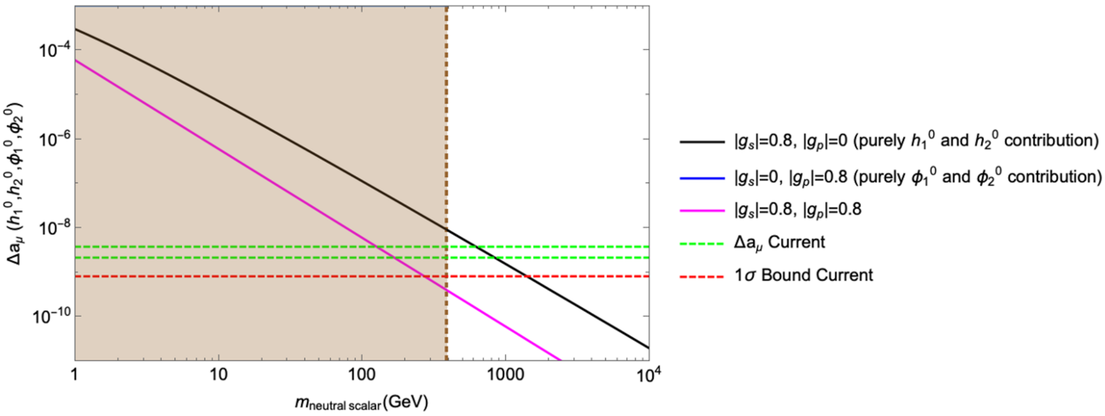

Figure 12. Plot showing the contributions of CP-even and CP-odd neutral scalars to Δa_μ. The black line represents purely scalar contribution coming from h_1^0, h_2^0, while the magenta line represents the combination of both scalar as well as pseudoscalar couplings. For purely pseudo-scalar coupling the contribution becomes negative (we have not shown it in this plot). Brown shaded region indicates the excluded mass range for neutral scalars.

the excluded range). Also from the results it can be concluded that singly charged scalars are not good candidates for explaining muon $(g-2)$ anomaly since they give negative and suppressed contribution.

Figure 12 shows the contributions coming from all the neutral scalars present in our model (all of them arising from bidoublet Φ), namely $h_1^0, h_2^0, \phi_1^0, \phi_2^0$. We have mentioned earlier that the contribution to muon anomaly coming from either pure scalar or pure pseudo-scalar or both can be easily derived from their couplings. In this case we can see that the neutral CP-even scalars h_1^0 and h_2^0 with mass around 500 GeV can explain the entire anomaly if we consider pure scalar couplings i.e., $g_s = 0.8$ and $g_p = 0$ for them (represented by black line). If we consider purely pseudo-scalar coupling; i.e. $g_s = 0$ and $g_p = 0.8$ then the contribution becomes negative (we have not plotted this contribution in figure 12). However if we take non-zero values for both the scalar and pseudo-scalar couplings, i.e. $g_s = 0.8$ and $g_p = 0.8$ (represented by the magenta line), then a neutral scalar with 150 GeV mass can address the anomaly since it is sensitive to the bounds on Δa_μ. But only the CP-even scalars can satisfy both muon anomaly as well as allowed mass range constraints for neutral scalars in one go (excluded mass range in the figure is indicated by brown shaded region). However considering massive CP even scalars with mass around $\mathcal{O}(10)$ TeV or higher, though saturate the allowed mass range constraint, fail to satisfy the current bound on muon anomaly. This can also be easily inferred from the plot. It is to be noted that neutral scalars are constrained by LEP searches for four-lepton contact interactions which requires $\frac{g}{M_\phi} < 2.5 \times 10^{-4}\,\text{GeV}^{-1}$ for $M_\phi > \sqrt{s}$ [129]. For our case $\frac{g}{M_\phi} = 1.6 \times 10^{-4}\,\text{GeV}^{-1}$ which clearly satisfies the LEP search bound.

Figure 13 shows the contribution coming from the new neutral vector boson $Z_{\mu\tau}$ in our model that comes from the $\text{U}(1)_{L_\mu - L_\tau}$ extension of LRSM. The plot shows that for coupling $g_{\mu\tau} = 8 \times 10^{-4}$, the neutral vector boson $Z_{\mu\tau}$ having mass nearly 150 MeV can

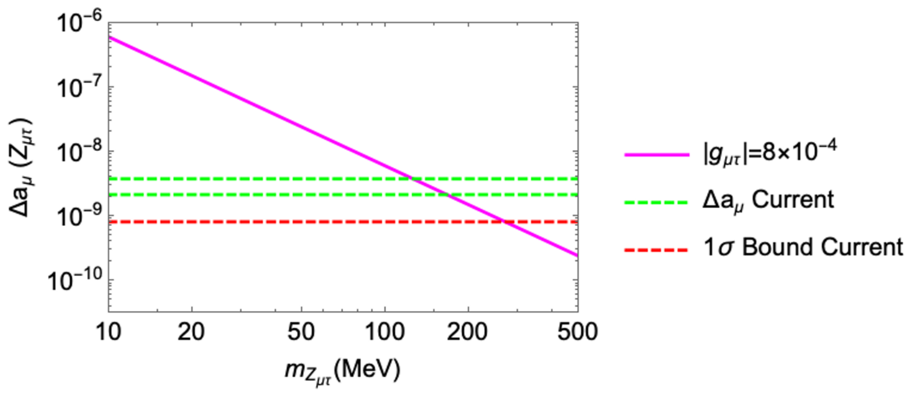

Figure 13. Plot showing the contribution coming from new light gauge boson $Z_{\mu\tau}$ *vs* mass of $Z_{\mu\tau}$. The magenta line shows the contribution of $Z_{\mu\tau}$ with coupling $g_{\mu\tau} = 8 \times 10^{-4}$ can address the anomaly with $Z_{\mu\tau}$ mass lying around 150 MeV.

address the entire anomaly (magenta line). The coupling strength ($g_{\mu\tau}$) of this vector boson is strongly constrained to be less than $\simeq 10^{-3}$ from the measurement of neutrino trident cross section by experiments like CHARM-II [131] and CCFR [132] while a mass of $\mathcal{O}(100\,\text{MeV})$ is allowed, and both of these are satisfied in our case.

In general, the individual contribution to muon anomaly arising from a mediating particle is related to its mass by the relation,

$$\Delta a_\mu \propto \frac{1}{m^2_{\text{mediator}}} \tag{5.2}$$

Thus all the contributions to Δa_μ become negligible in our model except those of W_L and $Z_{\mu\tau}$, if we consider heavy mass for the particles which are allowed by the collider experiments. The significant contribution comes from W_L if we consider large light-heavy neutrino mixing which is facilitated by inverse seesaw mechanism in the model. However, the light-heavy neutrino mixing parameter $V^{\nu\xi}$ can be constrained by other sectors like non-unitarity effects in experiments looking for lepton flavour violation and NSI effects at neutrino factory as discussed in section 3.1.

Figure 14 shows how the contribution of W_L to Δa_μ varies with different mixing values. Here we vary this mixing from 10^{-2} to 1. The magenta line represents the dependence of Δa_μ on light-heavy mixing and we find that $V^{\nu\xi}$ should be of $\mathcal{O}(0.3 - 1)$ in order to satisfy current bound on Δa_μ. A complete analysis of the results from the plots show that significant contributions to Δa_μ come from W_L and $Z_{\mu\tau}$ in the model whereas all other contributions are either negative, suppressed or ruled out by collider limits. However lighter neutral CP even scalars with mass around $0.5 - 2\,\text{TeV}$ can also be a good candidate to satisfy the entire current muon anomaly bound individually. Were the light-heavy neutrino mixing to be large in the inverse seesaw framework, W_L contribution could have accounted for the entire muon anomaly [1, 4] individually when the mixing ($V^{\nu\xi}$) is $\mathcal{O}(1)$. However the present bounds on the $|\eta|$ parameters of section 3.1 allow $\lesssim 0.3$ as the optimistic value of this parameter. If these bounds are confirmed this contribution is still capable

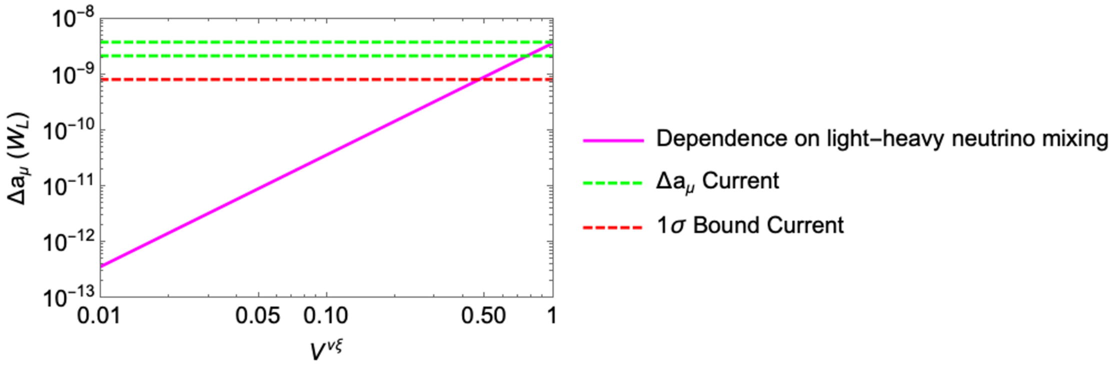

Figure 14. Plot showing the variation of Δa_μ coming from purely left-handed currents via W_L mediation vs. the light-heavy mixing parameter $V^{\nu\xi}$.

of explaining approximately 10% of the anomaly as can be seen from figure 14. Also for light extra neutral gauge boson contribution, we can easily infer that $m_{Z_{\mu\tau}} \sim 150\,\mathrm{MeV}$ and coupling $g_{\mu\tau} \sim 8 \times 10^{-4}$ can ameliorate the entire anomaly. Equally importantly we have thus established that in case natural values of the parameters of any one contribution are insufficient, the three together (i.e., contributions coming from CP even scalars h_1^0, h_2^0, singly-charged gauge boson W_L and new light neutral gauge boson $Z_{\mu\tau}$ mediation channels) have the potential to explain the entire anomaly within our ELISS scenario.

6 Conclusion

We have studied the $\mathrm{U}(1)_{L_\mu - L_\tau}$ extension of left-right symmetric model which can explain non-zero neutrino mass, mixing and muon anomalous magnetic moment within a single framework. Neutrino mass is generated in the model through inverse seesaw mechanism that allows large light-heavy neutrino mixing. We have discussed how the choice of scalars in various LRSM-SM symmetry breaking chains affect the generation of neutrino mass. We have calculated the individual contributions due to all the vector bosons and scalars present in the model to muon anomaly and found out that vector boson W_L with $\mathcal{O}(1)$ light-heavy neutrino mixing, the new light neutral vector boson $Z_{\mu\tau}$ as well as low-massive CP-even scalars are good candidates for explaining the entire anomaly. Although W_L's interaction with heavy right-handed neutrino, facilitated by inverse seesaw mechanism, becomes one of the significant contributions to the anomaly as this can account for upto 10% of the entire anomaly if one considers the constraints from NSI. Another major contribution comes from the new gauge boson $Z_{\mu\tau}$ which can explain the whole anomaly with mass 150 MeV and coupling $g_{\mu\tau} = 8 \times 10^{-4}$. The contributions coming from Z_R for different choices of couplings are negative whereas those of W_R are positive but invalid since it does not satisfy the allowed mass range. We have also briefly presented the comparative study of the effects between symmetric and asymmetric LRSM scenarios to muon anomaly estimation. Similarly the contributions arising from the charged as well as CP-odd neutral scalars are either suppressed or negative whereas CP-even neutral scalars can satisfy the entire muon anomaly bound for mass range $\sim 500\,\mathrm{GeV}$, but considering massive scalars with mass

$\sim \mathcal{O}(10)$ TeV or higher, we will get negligible contribution to muon anomaly due to neutral scalar mediation. We have also shown in plots how the contribution of each particle to Δa_μ varies with the mass of that particle for different choices of couplings. Overall we have found that inverse seesaw mechanism influences the results on muon anomaly to a large extent while explaining neutrino mass and mixing simultaneously in the model.

Acknowledgments

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JHEP09(2020)010