A simple model to explain the observed muon sector anomalies, small neutrino masses, baryon-genesis and dark-matter.

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Abstract

Since its inception, no decisive departure from the predictions of Standard Model (SM) has been reported. But recently various experiments have observed few hints of possible departure from SM predictions in lepton flavor universality observables such as $R_{K^*(1)}$, $P'_5$, muon $(g-2)$, $R(D^{(*)})$ etc. Many of these observable where deviation from SM in the range of $(2-4)\sigma$ were observed are related to muon ($\mu$) lepton. So these deviations may be some hint of a possible New Physics (NP) in the muon sector. In this work we extend the SM by introducing two SM singlet heavy charged leptons ($F_e$, $F_\mu$) whose left handed components are charged under a new $U(1)_F$ gauge symmetry, one color triplet lepto-quark ($\phi_Q$) doublet under $SU(2)_L$, one inert Higgs doublet ($\phi_l$), three very heavy Majorana neutrinos ($N_{iR}$) and one SM singlet scalar Dark Matter (DM) particle (S), all of which are odd under a $Z_2$ discrete symmetry. One more scalar ($\phi$) charged only under the $U(1)_F$ whose VEV give masses to the $U(1)_F$ gauge boson as well as the heavy leptons. With these new particles, we show that the observed anomalies in the muon sector as well as small neutrino masses, baryon-genesis and DM all can be explained with taking into account all the other experimental and theoretical constrains till date.

1 Introduction.

The Standard Model (SM) of particle physics turn out to be very simple but powerful mathematical construct that has stood unscathed from many experimental probes to find its loop holes for about forty years by now. Although SM itself has been verified by many experimental probes of its predictions, the

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neutrinos oscillation (the simplest way to interpret it, is to assume neutrinos have tiny but non-zero masses) and the dark-energy and dark-matter (DM) are clear indication of its incompleteness. But recently some intriguing anomalies has been reported by various experiments in muon (g-2) $^{[1]}$, $R_{K^{(*)}}^{[2][3][4]}$ and $P_{5}^{[3][6][7]}$ which may be indications of cracks in the SM. In general the flavor changing neutral current (FCNC) are sensitive to new-physics (NP) because SM is free of FCNC at tree level. One particular FCNC mode which has been well studied is $b \to s l\ell$, and it provoked tremendous interest in the particle physics community when LCHb reported anomalies in $B \to K^{*} \mu^{+} \mu^{-}$, $B_{s} \to \phi \mu^{+} \mu^{-}$ and $R_{K^{(*)}} = \frac{Br(B \to K^{(*)} \mu^{+} \mu^{-})}{Br(B \to K^{(*)} e^{+} e^{-})}^{[4]}$. Although the deviations in each individual modes are in the range of $(2.2 - 2.6)\sigma$, since all these mode are in the $b \to s \mu \mu$, the combine amounts to a significant deviation from SM prediction, though exact significant is still not properly settled yet, but some estimate put that to as high as $4\sigma^{[9]}$. A global fit to NP indicates that a NP contributions in Wilson coefficients $C_{9}$, $C_{9} = -C_{10}$, or $C_{9} = -C_{9}'$ with preference for large negative $C_{9}^{NP}$ at the level of $4-5\sigma$ than SM $^{[8][9][10]}$. In this work we propose a model where NP contribute to $b \to s \mu \mu$ via box diagram to generate a NP Wilson coefficient $C_{9}^{NP} = -C_{10}^{NP}$. The combine global fit in this case for the Wilson coefficients is given as $^{[9]}$

$$-0.81 \leq C_{9}^{NP} = -C_{10}^{NP} \leq -0.51 \text{ (at 1}\sigma)\text{.}$$

The most important constrains on the Wilson coefficients in these kind of models comes from the $B^{0} - \bar{B}^{0}$ mixing, $b \to s \gamma$ and $Br(B_{s}^{0} \to \mu \mu)$ where for the $B^{0} - \bar{B}^{0}$ mixing we have $^{[11]}$

$$C_{BB}(\mu_{H}) \epsilon [-2.1, 0.6] \times 10^{-5} TeV^{-2} \text{ (at 2}\sigma)\text{.}$$

where $\mu_{H} = 2m_{W}$ and for the $Br(B_{s}^{0} \to \mu^{+} \mu^{-})$ we have $^{[12]}$

$$Br(B_{s} \to \mu^{+} \mu^{-})_{Exp.} = 2.8 \begin{bmatrix} -0.7 \\ +0.7 \end{bmatrix} \times 10^{-9}$$

which is $1.2\sigma$ below the SM expectation of $Br(B_{s} \to \mu^{+} \mu^{-})_{SM} = (3.66 \pm 0.23) \times 10^{-9}^{[13]}$. For the $b \to s \gamma$, the constrain on $C_{7}^{NP}$ and $C_{8}^{NP}$ at $2\sigma$ turn out to be $^{[11]}$

$$-0.098 \leq C_{7}(\mu_{H}) + 0.24C_{8}(\mu_{H}) \leq 0.07 \text{ (at 2}\sigma),$$

where $\mu_{H}$ is take at $2m_{W}$. Then there is also the observed anomaly in the muon (g-2)

$$\delta a_{\mu} = a_{\mu}^{Exp.} - a_{\mu}^{SM} = 288(63)(49) \times 10^{-11},$$

which is at $3.6\sigma$ deviation from SM prediction $^{[1]}$. 

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2 Model details.

In this work we would like to propose a new-physics (NP) model which can explain the anomalies observed in $R_K$ and muon $(g-2)$ along with the loop generated neutrino masses and baryon-genesis via lepto-genesis. Also a viable scalar dark-matter (DM) is included. In line with the arguments given in [11][13], we introduce two heavy scalar $\phi_Q$ and $\phi_l$, where $\phi_Q$ is a lepto-quark, doublet under $SU(2)_L$ SM gauge group, and $\phi_l$ is the inert-doublet of inert two-Higgs-doublet model (IDM), both of which are odd under a discrete $Z_2$ transformation. Three heavy charged leptons $F_e$, $F_\mu$ and $F_\tau$ whose left-handed components are charged under a new $U(1)_F$ gauge symmetry, all are odd under the discrete $Z_2$ transformation where the subscripts on the heavy lepton symbols denote lepton numbers they carry. We also introduce a new scalar $\phi$ charged under the $U(1)_F$ which develops a non-zero VEV and gives masses to the new gauge boson $Z^F_\mu$ as well as new heavy leptons $F_e$, $F_\mu$ and $F_\tau$. In addition to the above new particles we also add three right-handed Majorana fermions $N_{iR}$ to generate neutrino masses at one loop as well as to generate the baryon asymmetry via lepton-genesis [15] and our model also include a scalar DM candidate S [16]. The new particles and their charge under the various transformations are tabulated in Table-1. With these new particles, as pointed out in [11][13], one loop box contribution to $b \rightarrow s \mu^+\mu^-$ can be generated. Due to the fact that left-handed $F_l$ being charged under the $U(1)_F$, only the right-handed components of the new particles and the left-handed components of SM fermions (which are both not charged under $U(1)_F$) can interact via Yukawa terms given as

$$L_{int} = \sum_i (y_i^Q \bar{Q}_i P_R F_i \phi_Q + y_i^L \bar{L}_i P_R F_i \phi_l) + h.c.$$  \hspace{1cm} (6)

Now the main constrains on the values of $Y_i$ and $n_i$ comes from the anomaly free conditions which gives

$$\sum_{i=e}^{\tau} Y_i^2 n_i = 0$$

$$\sum_{i=e}^{\tau} n_i^2 Y_i = 0$$

$$\sum_{i=e}^{\tau} n_i^3 = 0$$  \hspace{1cm} (7)

which are the anomaly free conditions coming from $U(1)_Y^2 U(1)_F$, $U(1)_F^2 U(1)_Y$ and $U(1)_F^3$ respectively and one more anomaly free condition due to gravity as $Gravity^2 U(1)_F$ which gives

$$\sum_{i=e}^{\tau} n_i = 0.$$  \hspace{1cm} (8)
| Particles | $SU(3)_c$ | $SU(2)_L$ | $U(1)_Y$ | $U(1)_F$ | $Z_2$ |
|-----------|-----------|-----------|-----------|-----------|-------|
| $\phi_Q$  | 3         | 2         | $7/6$     | 0         | -1    |
| $\phi_l$  | 1         | 2         | $1/2$     | 0         | -1    |
| $F_{iL}$  | 1         | 1         | $Y_i$     | $n_i$     | -1    |
| $F_{iR}$  | 1         | 1         | 0         | 0         | -1    |
| $N_{jR}$  | 1         | 1         | 0         | 0         | -1    |
| $S$       | 1         | 1         | 0         | 0         | -1    |
| $\phi$    | 1         | 1         | 0         | $n_\phi = n_\mu = -n_e$ | +1    |

Table 1: The charge assignments of new particles under the full gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y \times U(1)_F$. Where $i = e, \mu, \tau$ and $j = 1, 2, 3$.

The simplest non-trivial solution of the above four equations is given in [17] by setting $n_\tau = 0$, then $Y_\mu = -Y_e$ and $n_\mu = -n_e$ solves the above four equations with $Y_\tau$ a free parameter, in this work we do without requiring the existence of $F_\tau$. We take $Y_2 = Y_\mu = Q_{F_\mu} = -1$ then Yukawa terms given in Eqs.(6) for $F_1 = F_e$ is not allowed due to charge conservation and so no contribution to $b \to s e^+e^-$ from NP is expected, which is in line with the experimental findings that the NP is most likely in the muon sector instead of the electron sector [18]. And also, due to charge conservation and Lorenz invariance requirement, the $F_\epsilon$ will be a stable heavy charged lepton [2] whose mass (from the latest PDG [19] lower bound for heavy charged lepton mass) is $m_{F_e} \geq 102.6$ GeV but searches for long lived stable charged particles (in SUSY context) at LHC put the lower bound on heavy charged leptons as $m_{F_e} \geq 620$ GeV [20]. Now from the Eqs.(11) of [21] we have for the contribution of a neutral Higgs to the $\delta a_\mu$ coming from the Yukawa terms of Eqs.(6) is given as

$$
\delta a_\mu = \frac{m_\mu^2 y_\mu^2}{2 \times 16 \pi^2} \int_0^1 dx \left[ \frac{x^2 - x^3}{m_\mu^2 x^2 + (m_{F_e}^2 - m_\mu^2)x + m_{H_i}^2(1-x)} + \frac{x^2 - x^3}{m_\mu^2 x^2 + (m_{F_e}^2 - m_\mu^2)x + m_{A_i}^2(1-x)} \right]
$$

and in the limit $m_{A_i} \approx m_F \approx m_{H_i}^0 >> m_\mu$, $|m_F - m_{H_i}|$, we have

$$
\delta a_\mu \approx \frac{m_\mu^2}{12 \times 16 \pi^2} \left( \frac{y_\mu^2}{m_F^2} \right)^2,
$$

1. In general with $Y_\tau = 0$, the $F_\tau$ could also be a DM candidate, but in this work we will leave that possibility for a future pursuit and if we take $Y_\tau = -1$, then it can contribute to $b \to s \tau \tau$ as well as $(g-2)_\tau$.

2. $F_e$ can decay with introduction of (doubly charged under $U(1)_Y$ and singly charged under $U(1)_F$) scalars with Yukawa terms (at the order of LFV in charge leptons) such as $\bar{F}_{\mu R} \phi_{-n_e} F_{e L}$, with the $\phi_{-n_e}$ having large couplings to other exotic scalars, such that its present relic density would be negligible.
where if \( m_\mu > |m_F - m_{H_i}| \) then both \( F_\mu \) and \( H_i^0 \) will be stable, otherwise only the lighter of the two will be a stable particle. But in [32, 23] it has been shown that DM relic density contribution from scalar DM with Yukawa coupling in order unity (required for \( H_i^0 \) to explain the muon \((g-2)\) as shown below) is negligible, so in this work we introduce one more scalar singlet \( (S) \) to account the whole observed DM relic density. Experimentally the observed anomaly in the muon \((g-2)\) is given as

\[
\delta a_\mu = a_\mu^{Exp} - a_\mu^{SM} = 288(63)(49) \times 10^{-11}
\]

(amounting to about 3.6\( \sigma \) disagreement with SM prediction [1]. As first pointed out in [22], this discrepancy can be explained by a scalar and a heavy lepton \((F_\mu)\) propagating in the loop within 1\( \sigma \) of the experimental value with \( \frac{\delta a_\mu}{\sigma_{SM}} \approx 0.0188 \), for instance setting \( y_\mu = 3 \), \( m_{F_\mu} = 160 \) GeV, \( m_{H_i} = 150 \) GeV and \( m_{A_i} = 300 \) GeV in Eqs. (19), we get \( \delta a_\mu = 1.751 \times 10^{-9} \), which is within the 1.4\( \sigma \) of measured deviation. We will use \( \frac{\delta a_\mu}{\sigma_{SM}} \approx 0.0188 \) as a benchmark value in the following analysis. At \( Y_\mu = 3 \), \( m_{F_\mu} = 160 \) GeV, \( m_{H_i} = 150 \) GeV and \( m_{A_i} = 300 \) GeV, adopting formula II.39 of [23] and also see [17, 22], we get \( Br(Z \rightarrow \mu^+\mu^-)_{\text{triangle}} = 4.932 \times 10^{-6} \) compared to the experimental average of \( Br(Z \rightarrow \mu^+\mu^-)_{\text{exp.}} = (3.366 \pm 0.007)\% \) [1], the NP contribution via triangle loop is about an order of magnitude smaller than the errors in the most precise present experimental average from PDG. Collider signature of our model are similar to those given in [17].

2.1 Scalar Sector And DM.

The scalar potential can be written as

\[
V = \mu_1^2 |H|^2 + \mu_2^2 |\phi|^2 + \mu_3^2 |S|^2 + \mu_4^2 |\phi_Q|^2 + \mu_5^2 |\phi_l|^2 + \lambda_1 |H|^4 + \lambda_2 |\phi|^4 + \lambda_5 |S|^4 + \lambda_Q |\phi_Q|^4 + \lambda_l |\phi_l|^4 + \lambda_3 |H|^2 |\phi|^2 + \lambda_3QH |H|^2 |\phi_Q|^2 + \lambda_3SH |H|^2 |S|^2 + \lambda_3Q\phi |\phi_Q|^2 |\phi|^2 + \lambda_3\phi_l |\phi_l|^2 |\phi|^2 + \lambda_3Q|\phi|^2 |\phi_Q|^2 + \lambda_3\phi_l |\phi_l|^2 |\phi|^2 + \lambda_3Q|\phi_l|^2 |\phi_Q|^2 + \lambda_3S|\phi_l|^2 |S|^2 + \lambda_3S\phi |\phi_l|^2 |S|^2 + \lambda_4Q |H|^4 |\phi_Q|^2 + \lambda_4l |H|^4 |\phi_l|^2 + \lambda_4Q |\phi_l|^2 |H|^4 + \lambda_{5l}/2(|H|^4 |\phi_l|^2 + h.c.).
\]

(12)

where in the unitary gauges we have \( H = (0, \sqrt{2} (v + h))^T, \phi = \sqrt{2} (v_\phi + h_\phi), \phi_Q = (H_Q^{+5/3}, H_Q^{-2/3})^T \) and \( \phi_l = (H_l^+, \sqrt{2}(H_l^0 + iA_0^0))^T \) with \( S \) being a complex singlet scalar. Then the masses of the scalars are given by

\[
m_h^2 = \mu_1^2 + 3\lambda_1 v^2 + \lambda_3 v_\phi^2/2
\]

(13)

\[
m_h^2 = \mu_2^2 + 3\lambda_1 v_\phi^2 + \lambda_3 v^2/2
\]

(14)
\[ m_S^2 = \mu_S^2 + \lambda_{3SH} v^2 / 2 + \lambda_{3S\phi} v_\phi^2 / 2 \]  
\[ m_{H_Q^{+2/3}}^2 = m_{H_Q^{+5/3}}^2 + \lambda_{4Q} v^2 / 2 = \mu_Q^2 + \lambda_{3QH} v^2 / 2 + \lambda_{3\phi_Q} v_\phi^2 / 2 + \lambda_{4Q} v^2 / 2 \]  
\[ m_{H_H}^2 = \mu_l^2 + \lambda_{3H} v^2 / 2 + \lambda_{3\phi_l} v_\phi^2 / 2 \]  
\[ m_{H_H^0}^2 = \mu_l^2 + \lambda_{3H} v^2 / 2 + \lambda_{3\phi_l} v_\phi^2 / 2 + \lambda_{4H} v^2 / 2 + \lambda_{5H} v^2 / 2 \]  
\[ m_{A_Q}^2 = \mu_l^2 + \lambda_{3H} v^2 / 2 + \lambda_{3\phi_l} v_\phi^2 / 2 + \lambda_{4H} v^2 / 2 - \lambda_{5H} v^2 / 2. \]

As it is well known that the scalar singlet (S) stabilized by the \( Z_2 \) symmetry can be the DM candidate. If the couplings of the scalar singlet to other scalars are negligible compare to its coupling to Higgs, then the DM annihilation in this model is also dominated by the Higgs portal similar to the SM plus a scalar singlet models. In that case, recent XENON100 [24], SuperCDMS [25], LUX [26] and PandaX [27] measurements and limits on the invisible Higgs decay from LHC has ruled out the scalar singlet DM in the low mass region where the scalar singlet’s mass lies below Higgs mass except the narrow resonance region where it’s mass is close to \( m_h / 2 \). And recently the GAMBIT collaboration [16] has done a comprehensive fitting to all the measurements till date and found that the best fit point for the scalar singlet constituting the whole DM relic density is at \( \lambda_{3SH} = 2.9 \times 10^{-4} \) and \( m_S = 62.27 \) GeV in the low mass region and at \( \lambda_{3SH} = 3.1 \) and \( m_S = 9.79 \) TeV in the high mass region. But in our model there are two coupling constants not constrained by mass of the scalar singlet, i.e \( \lambda_{3S\phi_Q} \) and \( \lambda_{3S\phi_l} \), which can be adjusted such that even though direct, indirect and invisible Higgs decay etc. may require very small \( \lambda_{3SH} \), still \( \lambda_{3S\phi} \), \( \lambda_{3S\phi_Q} \) and \( \lambda_{3S\phi_l} \), will provide enough DM annihilation to avoid over abundance and so the entire mass range will be viable again (an exotic portal to the scalar singlet DM). Contributions from 2HDM to Peskin-Tekuchi \( \Delta T \) parameter (which is the most relevant parameter in 2HDM) can be made \( \Delta T \approx 0 \) by taking \( m_{H^\pm} \approx m_{H_H^0} \) [33].

### 3 Anomalies and bounds on Wilson coefficients.

In our model, which has same gauge group representation as the A-I of [11], the contribution from the NP to the observed anomalies in \( b \rightarrow s\mu\mu \) observables comes from the box loop, and in [11] the
authors have done a general analysis of such models. The contribution to the Wilson coefficients of $b \to s\mu\mu$ by NP box loop is given as \cite{11}

$$C_{NP}^9 = -C_{NP}^{10} = N \frac{Y_b Y_s^* |Y_\mu|^2}{2 \times 32\pi \alpha_{EM} m_{F_\mu}^2} [F(x_Q, x_{H_0}) + F(x_Q, x_{A_1})],$$

(20)

where $N^{-1} = \frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^*$ and

$$F(x, y) = \frac{1}{(1-x)(1-y)} + \frac{x^2 \ln[x]}{(1-x)^2(x-y)} + \frac{y^2 \ln[y]}{(1-y)^2(y-x)},$$

(21)

with $x_Q = \frac{m_{\phi_Q}^2}{m_{F_\mu}^2}$, $x_{H_0} = \frac{m_{H_0}^2}{m_{F_\mu}^2}$ and $x_{A_1} = \frac{m_{A_1}^2}{m_{F_\mu}^2}$. As mentioned in section 2 to explain the muon $(g-2)$ within $1\sigma$ of the experimental value, we need $Y_b Y_s^* \approx 0.0188$, so putting this value into Eqs.(20) and taking the benchmark values of the masses as $m_{\phi_Q} = 900$ GeV (which is about the present LHC lower bound \cite{29} \cite{30}), $m_{H_0} = 150$ GeV, $m_{A_1} = 300$ GeV and $m_{F_\mu} = 160$ GeV and setting $C_{NP}^9 = -C_{NP}^{10} = -0.66$, which is within the $1\sigma$ range of the present experimental bound given in Eqs.(11), we get $Y_b Y_s^* = -0.029$. At the above parameter values, we have $(C_9)^{NP}_{Penguin} \ll C_9^{NP} = -C_{10}^{NP}$. Using the benchmark values of masses and $Y_b Y_s^* = -0.029$ in

$$C_{BB} = \frac{(Y_b Y_s^*)^2}{128 \pi^2 m_{F_\mu}^2} F(x_Q, x_{Q})$$

(22)

we get $C_{BB} = 7.073 \times 10^{-7}$ TeV$^{-2}$ which is about an order of magnitude smaller than $2\sigma$ present experimental bound given in Eqs.(2) at $\mu_H = 2m_W$. Similarly with benchmark masses and $Y_b Y_s^* = -0.029$ we get

$$C_7(\mu_H) + 0.24 C_8(\mu_H) = -2.538 \times 10^{-3},$$

(23)

which is almost two-orders of magnitude smaller than the present $2\sigma$ experimental bound on this combination of Wilson coefficients coming from $b \to s \gamma$ data given in Eqs.(4), where

$$C_7 = \frac{NY_b Y_s^*}{2m_{F_\mu}^2} \left[ \frac{2}{3} F_7(x_Q) + \tilde{F}_7(x_Q) \right]$$

(24)

and

$$C_8 = \frac{NY_b Y_s^*}{2m_{F_\mu}^2} [F_7(x_Q)]$$

(25)

with $F_7(x) = \frac{x^3 - 6x^2 + 6\ln(x) + 3x + 2}{12(x-1)^4}$ and $\tilde{F}_7(x) = \frac{1}{x} F_7(x^{-1})$. Now another key observable that put very stringent constrain comes from the measurement of $Br(B_s \to \mu^+ \mu^-)_{Exp.} = 2.8 \begin{bmatrix} +0.7 \\ -0.6 \end{bmatrix} \times 10^{-9}$ which
is about 1.2σ below the SM prediction of $Br(B_s \to \mu^+\mu^-)_{SM} = (3.66 \pm 0.23) \times 10^{-9}$ \cite{28}. The decay $B_s \to \mu^+\mu^-$ can be expressed as

$$Br(B_s \to \mu^+\mu^-)_{\text{eff.}} = \frac{G_F^2 \alpha EM}{16\pi^3} |V_{tb}V_{ts}^\ast|^2 |C_{10}^{\text{eff.}}|^2 m_{B_s}^2 f_{B_s}^2 (1 + \mathcal{O}(m_{\mu}^2/m_{B_s}^2)),$$

(26)

where $C_{10}^{\text{eff.}} = C_{10}^{SM} + C_{10}^{NP}$ with $C_{9,10}^{SM} = (4.07, -4.31)$ \cite{9} and $C_{10}^{NP} = +0.66$ as our benchmark value, we get $Br(B_s \to \mu^+\mu^-)_{\text{eff.}} = 2.63 \times 10^{-9}$ which is well within the 1σ of the measured value. The bound coming from $B \to K^{(*)}\nu\nu$ is much weaker than the experimental bounds from $B \to K^{(*)}\mu^+\mu^-$, so we can ignore constrain from this mode \cite{11}.

4 Loop generation of neutrino masses and Baryon-genesis.

With presence of $N_{jR}$ we can have Yukawa terms such as

$$\mathcal{L}_Y = \sum_{i,j=1}^{3} h_{ij} \bar{L}_i \sigma_2 \phi_t N_{jR} + h.c.,$$

(27)

which is well known to give Majorana neutrino mass term $M_{\alpha\beta} \bar{\nu}_\alpha \nu_\beta + h.c.$ at one loop level via the scotogenic mechanism given as \cite{15}

$$M_{\alpha\beta} = \sum_i \frac{h_{\alpha i} h_{\beta i} M_i}{16 \pi^2} \left[ \frac{m_{H_0}^2}{m_{H_0}^2 - M_i^2} \ln \frac{m_{H_0}^2}{M_i^2} - \frac{m_{A_0}^2}{m_{A_0}^2 - M_i^2} \ln \frac{m_{A_0}^2}{M_i^2} \right],$$

(28)

where $m_{H_0}$ and $m_{A_0}$ are masses of the $H_0$ and $A_0$ respectively and $M_i$’s are the very heavy Majorana masses of the $N_{iR}$ neutrinos. In our benchmark values where $m_{H_0} = 150$ GeV and $m_{F_\mu} = 160$ GeV with assuming $m_{H_0} < m_{A_0} = 300$ GeV, the $H_0$ will be a DM candidate but due to requirement of large Yukawa coupling between $H_0$, $F_\mu$ and $\mu$ to explain the muon (g-2) data, as pointed out in \cite{23}, for such large Yukawa couplings the contribution to the present relic density of the DM by $H_0$ will be negligibly small. Now if the the terms inside the bracket in Eqs.\cite{28} is at the order of $\frac{\bar{\nu}_0^2}{M_i^2}$, then the lightest of the right handed Majorana neutrino need only be heavier than about $2.6 \times 10^7$ to make baryon-genesis via lepto-genesis possible, where $\bar{\nu}_0 \approx 1.6$ TeV. In the process of excess baryon generation from lepton excess generated by the CP violation introduced due to interference between the tree level and one loop level terms contributions to the decay $N_{1R} \to l_i + H_i^+$ \cite{31}, at the time of electro-weak epoch the baryon number excess and lepton number excess are related by \cite{31}

$$\delta B = - (\delta L)/2,$$

(29)

where $N_{1R}$ being the lightest of the three heavy Majorana neutrinos. A Yukawa couplings of order $|h_{11}|^2 \approx 10^{-7}$ will be able to generate required baryon number excess and neutrino masses of order
O(0.04) eV which is close to the latest experimental bound on largest of neutrino mass difference from neutrino mixing measurements of $|\Delta m_{32}| \approx 0.05$ eV [1], for more details see also [15] [31].

5 Conclusions.

In this work we have proposed a simple model which can explain the observed muon related anomalies along with small neutrino masses and baryon genesis. We have introduced one leptoquark ($\phi_Q$ which is triplet under $SU(3)_c$) and one inert Higgs doublet ($\phi_l$), both are odd under a $Z_2$ and doublet under $SU(2)_L$, at least two $SU(2)_L$ singlet heavy leptons $F_e$ and $F_\mu$, both odd under a $Z_2$ and whose left handed components are charged under a new $U(1)_F$ gauge symmetry. One $Z_2$ even scalar singlet under the SM gauge groups but charged under the new $U(1)_F$ gauge symmetry whose VEV gives masses to the new heavy leptons and $U(1)_F$ gauge boson. Also we added three very heavy right handed Majorana neutrinos odd under the $Z_2$ to generate neutrino masses at one loop via the scotogenic mechanism as well as to account for the observed baryon asymmetry via lepto-genesis. One more scalar $S$, odd under the $Z_2$, to account for the Dark Matter (DM) in the universe is also introduced.

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