

Non-universal electroweak extensions of the standard model and the scotogenic models.

Dario J. Bandon^{*a}, Diego A. Restrepo^{†a}, William A. Ponce^{‡a}, Eduardo Rojas^{§b},

^aInstituto de Física, Universidad de Antioquia,
Medellin, Antioquia

^bDepartamento de Física, Universidad de Nariño,
A.A. 1175, San Juan de Pasto, Colombia

Abstract

In order to analyze some low energy experimental anomalies, we charge with a non-universal $U(1)'$ gauge symmetry the standard model fermions, taking as a starting point the well-known scotogenic model. In order to have non-trivial solutions to the anomalies and the Yukawa constraints, we add three neutral singlet Dirac fermions. We have found two possible non-universal solutions which, as a matter of principle, are suitable to analyze family-dependent experimental anomalies.

1 Introduction

In the SM the couplings of the photon and the Z boson to the standard model (SM) fermions are universal, that is to say, the couplings of these bosons to the corresponding fermions in each family are the same, having as a consequence that couplings of these bosons to the fermions of the SM continue to be diagonal after the mixing between fermions [1]. However, there are several observables at low energies for which their experimental values get apart from those predicted by the SM. In many of these cases the best models to fit these anomalies are non-universal electroweak extensions to the SM; being this feature the main motivation for the present work. Some of the observables that show deviations from the SM are: the anomalous magnetic moment of the muon [2], the beryllium anomaly [3], MiniBooNE [4], Gallium solar neutrino experiments GALLEX [5, 6, 7] and SAGE [8, 9, 10, 11], NuTeV [12], LSND [13], and the Reactor anomaly [14]. There is also an increasing interest in a number of anomalies in semileptonic B decays reported by the LHCb collaboration and other experiments [15, 16, 17, 18, 19, 20, 21].

On the other hand, in neutrino physics the scotogenic models are compelling from a theoretical point of view, because they manage to link the dark matter and the mechanism to generate the neutrino masses [22]. Following well-established methods to induce residual symmetries from $U(1)$ -invariant theories [23, 24], within the framework of higher gauge symmetries, the dark matter stability can be explained through residual symmetries at low energies [25, 26, 27, 28]. In these models the neutrino masses are generated radiatively via the effective the Weinberg operator, in

*e-mail: dario.bandon@udea.edu.co

†e-mail: restrepo@udea.edu.co

‡e-mail: willian.ponce@udea.edu.co

§e-mail: eduro4000@gmail.com

such a way that the mechanism of generation of masses of neutrinos and the dark matter are related to each other. As it has been proposed in some recent works, the stability of dark matter can be explained by residual symmetries of gauge groups; avoiding the need to impose ad hoc symmetries. This procedure has been implemented in several models and in particular in some scotogenic models.

These results motivate the study of non-universal models and in particular those models that adjust some of these anomalies simultaneously. Motivated by this phenomenology, and bearing in mind that many of these anomalies suggest non-universal models, we extend the electroweak sector of the SM with an additional symmetry $U(1)$ with non-universal charges to the fermions of the SM. Additionally, we want our model to be able to explain masses of neutrinos and dark matter stability. As we mentioned earlier the scotogenic paradigm is the preferred scenario in these cases.

The manuscript is organized as follows: in section 2 write the anomaly equations and the restrictions coming from the Yukawa terms and the 1-loop neutrino self-energy diagram. In section 3 we will write our results.

2 Anomaly and Yukawa constraints

Following ref [29] we generalize that model by allowing Z' family-dependent charges. In table 1 l_{L_i} , e_{R_i} , q_{L_i} , u_{R_i} and d_{R_i} represent the SM fields associated with the i -th family. In order to have non-trivial solutions our model contains, two scalar doublets $H_{1,2}$ (as in the original model) and three scalar doublets Φ_i one for each family (two additional fields when comparing with the original reference). We also have three heavy fermion fields N_i , which have vector couplings under the extra $U'(1)$ abelian gauge symmetry and singlets under the SM gauge group.

| Field | Spin | $SU(3)_c$ | $SU(2)_L$ | $U(1)_Y$ | $U(1)'$ |
|--------------|---------------|-----------|-----------|----------------|--------------|
| l_{Lr} | $\frac{1}{2}$ | 1 | 2 | $-\frac{1}{2}$ | l_r |
| e_{Rr} | $\frac{1}{2}$ | 1 | 1 | -1 | e_r |
| q_{Rr} | $\frac{1}{2}$ | 3 | 2 | $\frac{1}{6}$ | q_r |
| u_{Rr} | $\frac{1}{2}$ | 3 | 1 | $\frac{2}{3}$ | u_r |
| d_{Rr} | $\frac{1}{2}$ | 3 | 1 | $-\frac{1}{3}$ | d_r |
| $N_{(L,R)r}$ | $\frac{1}{2}$ | 1 | 1 | 0 | n_r |
| Φ_r | 0 | 1 | 2 | $\frac{1}{2}$ | ϕ_r |
| $H_{1,2}$ | 0 | 1 | 2 | $\frac{1}{2}$ | $\eta_{1,2}$ |

Table 1: Particle content and quantum numbers under the $SU(3)_C \otimes SU(2)_L \otimes U(1) \otimes U(1)'$ gauge group. The indices $r = 1, 2, 3$ run over the three families

For the $SU(2)_L \otimes U(1) \otimes U(1)'$ symmetry the non-trivial anomaly equations are:

$$\begin{aligned}
[SU(2)]^2 U(1)' : 0 &= \Sigma q + \frac{1}{3} \Sigma l, \\
[SU(3)]^2 U(1)' : 0 &= 2\Sigma q - \Sigma u - \Sigma d, \\
[\text{grav}]^2 U(1)' : 0 &= 6\Sigma q - 3(\Sigma u + \Sigma d) + 2\Sigma l - \Sigma n - \Sigma e + \Sigma N \\
[U(1)]^2 U(1)' : 0 &= \frac{1}{3} \Sigma q - \frac{8}{3} \Sigma u - \frac{2}{3} \Sigma d + \Sigma l - 2\Sigma e \\
U(1)[U(1)']^2 : 0 &= \Sigma q^2 - 2\Sigma u^2 + \Sigma d^2 - \Sigma l^2 + \Sigma e^2, \\
[U(1)']^3 : 0 &= 6\Sigma q^3 - 3(\Sigma u^3 + \Sigma d^3) + 2\Sigma l^3 - \Sigma n^3 - \Sigma e^3 + \Sigma N^3.
\end{aligned} \tag{1}$$

where $\Sigma f = f_1 + f_2 + f_3$. We have also considered the constraints coming from Yukawa interaction terms

$$\mathcal{L}_Y \supset \sum_{r=1,2,3} \bar{l}_r \Phi_r e_r + \bar{q}_r \tilde{\Phi}_r u_r + \bar{q}_r \Phi_r d_r + \mu_{N_r}^2 \bar{N}_{R_r} N_{L_r} + \lambda (\Phi_1^\dagger H_1) (\Phi_1^\dagger H_2) + \text{H.C}$$

In addition to this Lagrangian, by following [29], we also added the restrictions that come from the terms $\sum_{r=j,k} \nu_{Lr} \eta_1^0 \bar{N}_{Rr}$, $\sum_{r=j,k} N_{rL} \nu_{Lr} \eta_2^0$, where, $H_{1,2}^T = (\eta_{1,2}^+, \eta_{1,2}^0)$. From these terms we obtain

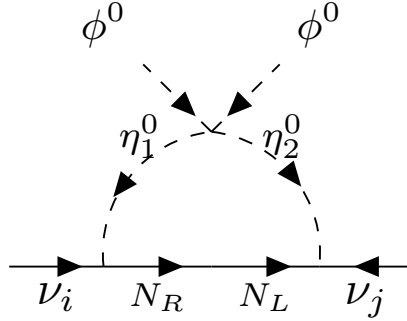


Figure 1: One-loop neutrino self-energy.

the constraints

$$\begin{aligned}
e_r - l_r + \phi_r &= 0, & d_r - q_r + \phi_r &= 0, & u_r - q_r - \phi_r &= 0, \\
l_j(l_k) + \eta_1 - n_j(n_k) &= 0, & l_j(l_k) + \eta_2 + n_j(n_k) &= 0, & +\eta_1 + \eta_2 - 2\phi_{j,k} &= 0
\end{aligned} \tag{2}$$

There are several options for the choice of the indices; in general, the triplet (ijk) is a permutation of (123) and $r = 1, 2, 3$. It is possible to combine the indices j, k in three different ways, i.e., $(1,2)$, $(1,3)$ or $(2,3)$. Each of these options represents a different model.

3 Results and conclusions

The solution of the equations (1) and (2) are shown in tables 2 and 3. The couplings to SM leptons and the scalar fields Φ_r in table 2, are the same for the three families; hence, for this solution the three fields Φ_r can be reduced to only one. That is interesting since in this case. the number of fields

| f | $\epsilon^{Z'}(f)$ | | |
|-------|--------------------|----------|-----------------------|
| l_r | $-\Sigma q$ | ϕ_r | $+\Sigma q$ |
| e_r | $-2\Sigma q$ | n^j | $+\eta_1 - \Sigma q$ |
| u_r | $+q_r + \Sigma q$ | n^k | $+\eta_1 - \Sigma q$ |
| d_r | $+q_r - \Sigma q$ | η_2 | $-\eta_1 + 2\Sigma q$ |

Table 2: Universal solution in the leptonic sector for the Z' charges in the equations (1) and (2). The free parameters are: the left-handed couplings to the quarks q_k , the heavy-vector field n_i , and the coupling to the H_1 field η_1 . The integers (ijk) are a permutation of (123) and they correspond to the family number.

| f | $\epsilon^{Z'}(f)$ | | | | | |
|----------|--------------------------|----------|---------------------------------------|----------|---------------------------------------|--|
| l_i | $-3q_i$ | l_j | $-\frac{3}{2}(q_j + q_k)$ | l_k | $-\frac{3}{2}(q_j + q_k)$ | |
| e_i | $-6q_i$ | e_j | $-3(q_j + q_k)$ | e_k | $-3(q_j + q_k)$ | |
| u_i | $+4q_i$ | u_j | $+\frac{1}{2}(5q_j + 3q_k)$ | u_k | $+\frac{1}{2}(3q_j + 5q_k)$ | |
| d_i | $-2q_i$ | d_j | $-\frac{1}{2}(q_j + 3q_k)$ | d_k | $-\frac{1}{2}(3q_j + q_k)$ | |
| ϕ_i | $+3q_i$ | ϕ_j | $+\frac{3}{2}(q_j + q_k)$ | ϕ_j | $+\frac{3}{2}(q_j + q_k)$ | |
| η_2 | $-\eta_1 + 3(q_j + q_k)$ | n_j | $+\frac{1}{2}(2\eta_1 - 3q_j - 3q_k)$ | n_k | $+\frac{3}{2}(2\eta_1 - 3q_j - 3q_k)$ | |

Table 3: Solution for the Z' charges in the equations (1) and (2). The free parameters are: the left-handed couplings to the quarks q_k , the heavy-vector field n_i , and the coupling to the H_1 field η_1 . The integers (ijk) are a permutation of (123) and they correspond to the family number.

is identical to the original model [29] even though the model is not universal. For $q_1 = q_2 = q_3 = 0$ and $\eta_1 = 1$ from this solution it is possible to obtain the original model [29]. The solution in the table 3 is more interesting for flavor physics, due to the fact that there is less dependence between the Z' charges of the particles in different families. By choosing $(i, j, k) = (1, 2, 3)$ and $q_1 = q_2$, it is possible to obtain a model with generation dependent charges in the lepton sector. By a convenient mixing matrix for the right-handed fields, it is possible to obtain a model without flavor changing neutral currents. This model is also important since it represents a realization of a non-universal Z' model without right-handed neutrinos with allowed one-loop contributions to the neutrino masses. As explained in reference [29] the dark matter candidate is either, the real part of $\eta_{1,2}^0$ or one of the exotic fermions N_i .

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