# Radiative Left-Right symmetry breaking from flavour enhanced trinification

**Jonas Wessén**[1]

J. E. Camargo-Molina[1] A. P. Morais[1,2] R. Pasechnik[1]

[1]Dept. of Astronomy and Theoretical Physics
Lund University, Lund, Sweden

[2]Center for Research and Development in Mathematics and Applications (CIDMA)
Aveiro University, Aveiro, Portugal

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## Overview:

- The Standard Model (SM) is experimentally amazingly succesful, but contains theoretical issues, e.g.:
  - $\sim 20$ free parameters is not so elegant!
  - Why 3 families with large mass hierarchies? (e.g. $m_u/m_t \sim 10^{-5}$, $m_e/m_\tau \sim 10^{-4}$)
- We postulate a high scale GUT based on the trinification gauge group $[\mathrm{SU}(3)]^3 \times \mathbb{Z}_3$, which is a maximal subgroup of $\mathrm{E}_6$.
- Scalars and fermions form bi-triplet gauge representations, which fit into the **27** rep of $\mathrm{E}_6$.
- A global family symmetry $\mathrm{SU}(3)_\mathrm{F}$ is imposed which can be interpreted as originating from $\mathrm{E}_8 \to \mathrm{E}_6 \times \mathrm{SU}(3)$. This leads to a very constrained model with 12 free parameters!
- Trinification is spontaneously broken at tree-level to a smaller Left-Right symmetric (LR) gauge group.
- Integrating out the heavy fields leads to an effective LR model. Upon RG running down towards the electro-weak scale, the gauge group is broken to the SM gauge group due a sign flip in a squared scalar mass.
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Gauge | Global

$$[\mathrm{SU(3)_L \times SU(3)_R \times SU(3)_C}] \times \mathbb{Z}_3 \times \{\mathrm{SU(3)_F \times U(1)_A \times U(1)_B}\}$$

$\downarrow$ (tree-level)

$$\mathrm{SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{L+R}} \times \{\mathrm{SU(2)_F \times U(1)_X \times U(1)_Z \times U(1)_B}\}$$

$\downarrow$ (radiative)

$$\mathrm{SU(3)_C \times SU(2)_L \times U(1)_Y} \times \{\mathrm{U(1)_D \times U(1)_E \times U(1)_G \times U(1)_B}\}$$

$\downarrow$ (radiative?)

$$\mathrm{SU(3)_C \times U(1)_{E.M.}} \times \{(\,?\,) \times \mathrm{U(1)_B}\}$$

| | $\mathrm{SU(3)_L}$ | $\mathrm{SU(3)_R}$ | $\mathrm{SU(3)_C}$ | $(\mathrm{SU(3)_F})$ |
|---|---|---|---|---|
| Left-handed Weyl Fermions | | | | |
| $L$ | $\mathbf{3}$ | $\bar{\mathbf{3}}$ | $\mathbf{1}$ | $\mathbf{3}$ |
| $Q_\mathrm{L}$ | $\bar{\mathbf{3}}$ | $\mathbf{1}$ | $\mathbf{3}$ | $\mathbf{3}$ |
| $Q_\mathrm{R}$ | $\mathbf{1}$ | $\mathbf{3}$ | $\bar{\mathbf{3}}$ | $\mathbf{3}$ |
| Scalars | | | | |
| $\tilde{L}$ | $\mathbf{3}$ | $\bar{\mathbf{3}}$ | $\mathbf{1}$ | $\mathbf{3}$ |
| $\tilde{Q}_\mathrm{L}$ | $\bar{\mathbf{3}}$ | $\mathbf{1}$ | $\mathbf{3}$ | $\mathbf{3}$ |
| $\tilde{Q}_\mathrm{R}$ | $\mathbf{1}$ | $\mathbf{3}$ | $\bar{\mathbf{3}}$ | $\mathbf{3}$ |
| Gauge Bosons | | | | |
| $G_\mathrm{L}$ | $\mathbf{8}$ | $\mathbf{1}$ | $\mathbf{1}$ | $\mathbf{1}$ |
| $G_\mathrm{R}$ | $\mathbf{1}$ | $\mathbf{8}$ | $\mathbf{1}$ | $\mathbf{1}$ |
| $G_\mathrm{C}$ | $\mathbf{1}$ | $\mathbf{1}$ | $\mathbf{8}$ | $\mathbf{1}$ |

The $\mathbb{Z}_3$ refers to a cyclic permutation of the fields $\{L, Q_\mathrm{L}, Q_\mathrm{R}\}$, $\{\tilde{L}, \tilde{Q}_\mathrm{L}, \tilde{Q}_\mathrm{R}\}$ and $\{G_\mathrm{L}, G_\mathrm{R}, G_\mathrm{C}\}$, and enforces gauge coupling unification.

**Unification of Yukawa couplings**

Before SSB, all fermions are massless and only one Yukawa term is allowed:

$$\mathcal{L}_\mathrm{F} = -y\, \epsilon_{ijk}\, (\tilde{L}^i)^l{}_r\, ({Q_\mathrm{L}}^j)^c{}_l\, ({Q_\mathrm{R}}^k)^r{}_c + \mathrm{c.c.} + (\mathbb{Z}_3 \text{ permutations}).$$

Scalar potential: $V_\mathrm{S} = V_1 + V_2 + V_3$ with

$$\begin{aligned}
V_1 = & -\mu^2 (\tilde{L}^i)^l{}_r (\tilde{L}^*_i)_l{}^r + \lambda_1 \left[ (\tilde{L}^i)^l{}_r (\tilde{L}^*_i)_l{}^r \right]^2 \\
& + \lambda_2 (\tilde{L}^i)^l{}_r (\tilde{L}^j)^{l'}{}_{r'} (\tilde{L}^*_j)_l{}^r (\tilde{L}^*_i)_{l'}{}^{r'} \\
& + \lambda_3 (\tilde{L}^i)^l{}_r (\tilde{L}^j)^{l'}{}_{r'} (\tilde{L}^*_i)_l{}^{r'} (\tilde{L}^*_j)_{l'}{}^r \\
& + \lambda_4 (\tilde{L}^i)^l{}_r (\tilde{L}^j)^{l'}{}_{r'} (\tilde{L}^*_j)_l{}^{r'} (\tilde{L}^*_i)_{l'}{}^r \\
& + \ (\mathbb{Z}_3 \text{ permutations}), \\
V_2 = & \ \alpha_1 (\tilde{L}^i)^l{}_r (\tilde{L}^*_i)_l{}^r (\tilde{Q}_\mathrm{L}{}^j)^c{}_{l'} (\tilde{Q}^*_{\mathrm{L}j})_c{}^{l'} \\
& + \alpha_2 (\tilde{L}^i)^l{}_r (\tilde{L}^*_j)_l{}^r (\tilde{Q}_\mathrm{L}{}^j)^c{}_{l'} (\tilde{Q}^*_{\mathrm{L}i})_c{}^{l'} \\
& + \alpha_3 (\tilde{L}^i)^l{}_r (\tilde{L}^*_i)_{l'}{}^r (\tilde{Q}_\mathrm{L}{}^j)^c{}_l (\tilde{Q}^*_{\mathrm{L}j})_c{}^{l'} \\
& + \alpha_4 (\tilde{L}^i)^l{}_r (\tilde{L}^*_j)_{l'}{}^r (\tilde{Q}_\mathrm{L}{}^j)^c{}_l (\tilde{Q}^*_{\mathrm{L}i})_c{}^{l'} \\
& + \ (\mathbb{Z}_3 \text{ permutations}), \\
V_3 = & \ \gamma\, \epsilon_{ijk} (\tilde{L}^i)^l{}_r (\tilde{Q}_\mathrm{L}{}^j)^c{}_l (\tilde{Q}_\mathrm{R}{}^k)^r{}_c + \mathrm{c.c.}
\end{aligned}$$

$V_{\rm S}$ has a minimum where

$$\langle (\tilde{L}^i)^l{}_r \rangle = \delta^i_3 \delta^l_3 \delta^3_r \frac{v_3}{\sqrt{2}} = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & \frac{v_3}{\sqrt{2}} \end{pmatrix}^{i=3}$$

This is equivalent to putting the VEV in either $\tilde{Q}_{\rm L}$ or $\tilde{Q}_{\rm R}$ due to the $\mathbb{Z}_3$ symmetry.

**Scalar mass eigenstates:**

| Fields | $(\text{Mass})^2$ | Comment |
|---|---|---|
| $(\tilde{L}^l)^L{}_R$ | $-(\lambda_2+\lambda_3+\lambda_4)\,v_3^2$ | |
| $(\tilde{L}^l)^3{}_R$ | $-(\lambda_2+\lambda_3)\,v_3^2$ | |
| $(\tilde{L}^3)^L{}_R$ | $-(\lambda_3+\lambda_4)\,v_3^2$ | |
| $(\tilde{L}^l)^L{}_3$ | $-(\lambda_2+\lambda_4)\,v_3^2$ | |
| $\mathrm{Re}[(\tilde{L}^3)^3{}_3]$ | $2(\lambda_1+\lambda_2+\lambda_3+\lambda_4)\,v_3^2$ | |
| $\mathrm{Im}[(\tilde{L}^3)^3{}_3]$ | 0 | Gauge Goldstone |
| $(\tilde{L}^3)^L{}_3$ | 0 | Gauge Goldstone |
| $(\tilde{L}^3)^3{}_R$ | 0 | Gauge Goldstone |
| $(\tilde{L}^l)^3{}_3$ | 0 | Global Goldstone |

(+ All $\tilde{Q}_{\mathrm{L,R}}$ get masses of $\mathcal{O}(v_3)$)

**Fermion mass eigenstates:**

| Fields | $(\text{Mass})^2$ | Comment |
|---|---|---|
| $\begin{pmatrix} ({Q_{\mathrm{L}}}^{I})^{c}{}_{3} \\ \epsilon^{IJ}(Q^{\dagger}_{\mathrm{R}\,J})_{3}{}^{c} \end{pmatrix}$ | $\frac{1}{2}y^2 v_3^2$ | $\mathrm{SU}(2)_{\mathrm{F}}$ doublet |

Table: Among the fermions, 1 exotic flavour doublet heavy quark gets a (Dirac) mass. All others are massless at this stage.

## The low-scale effective Left-Right symmetric theory

For $\mathrm{SU}(2)_\mathrm{R} \times \mathrm{U}(1)_\mathrm{L+R} \to \mathrm{U}(1)_\mathrm{Y}$ breaking we need a vev in

$$\tilde{R}^{l}{}_{R} \equiv (\tilde{L}^{l})^{3}{}_{R}.$$

Also need at least one $\mathrm{SU}(2)_\mathrm{L}$ doublet to trigger EWSB $\mathrm{SU}(2)_\mathrm{L} \times \mathrm{U}(1)_\mathrm{Y} \to \mathrm{U}(1)_\mathrm{E.M.}$, e.g.

$$\tilde{h}^{L}{}_{R} \equiv (\tilde{L}^{3})^{L}{}_{R}.$$

We define

$$\begin{aligned} m_R^2 &= -(\lambda_3 + \lambda_2)\, v_3^2 &\equiv \epsilon\, v_3^2, \\ m_h^2 &= -(\lambda_3 + \lambda_4)\, v_3^2 &\equiv \delta\, v_3^2. \end{aligned}$$

and integrate out all heavy states except $\tilde{h}^{L}{}_{R}$ and $\tilde{R}^{l}{}_{R}$, assuming $\delta, \epsilon \ll 1$.

| | $\mathrm{SU(2)_L}$ | $\mathrm{SU(2)_R}$ | $\mathrm{SU(3)_C}$ | $\mathrm{U(1)_{L+R}}$ | $(\mathrm{SU(2)_F})$ | $(X, Z, B)$ |
|---|---|---|---|---|---|---|
| Scalars | | | | | | |
| $\tilde{h}$ | $\mathbf{2}$ | $\bar{\mathbf{2}}$ | $\mathbf{1}$ | 0 | $\mathbf{1}$ | $(+2, 0, 0)$ |
| $\tilde{R}$ | $\mathbf{1}$ | $\bar{\mathbf{2}}$ | $\mathbf{1}$ | 1 | $\mathbf{2}$ | $(-1, 1, 0)$ |
| $\tilde{\Phi}$ | $\mathbf{1}$ | $\mathbf{1}$ | $\mathbf{1}$ | 0 | $\mathbf{2}$ | $(-2, 1, 0)$ |

$$
\begin{aligned}
V_S^{LR} = {} & m_h^2|\tilde{h}|^2 + m_R^2|\tilde{R}|^2 + m_{\tilde{\Phi}}^2|\tilde{\Phi}|^2 \\
& + \lambda_a|\tilde{h}|^4 + \lambda_b|\tilde{R}|^4 + \lambda_c|\tilde{\Phi}|^4 + \lambda_d|\tilde{h}|^2|\tilde{\Phi}|^2 + \lambda_e|\tilde{R}|^2|\tilde{\Phi}|^2 + \lambda_f|\tilde{R}|^2|\tilde{h}|^2 \\
& + \lambda_g(\tilde{R}^l)_{R_1}(\tilde{R}^*_l)^{R_2}(\tilde{h})^L{}_{R_1'}(\tilde{h}^*)_L{}^{R_2'} \quad \epsilon^{R_1 R_1'}\,\epsilon_{R_2 R_2'} \\
& + \lambda_h(\tilde{R}^{l_1})_R(\tilde{R}^*_{J_1})^R\,\tilde{\Phi}^{l_2}\tilde{\Phi}^*_{J_2} \quad \epsilon_{l_1 l_2}\,\epsilon^{J_1 J_2} \\
& + \lambda_i(\tilde{R}^{l_1})_{R_1}(\tilde{R}^{l_2})_{R_1'}(\tilde{R}^*_{J_1})^{R_2}(\tilde{R}^*_{J_2})^{R_2'} \quad \epsilon_{l_1 l_2}\,\epsilon^{J_1 J_2}\,\epsilon_{R_2 R_2'}\,\epsilon^{R_1 R_1'} \\
& + \lambda_j(\tilde{h})^{L_1}{}_{R_1}(\tilde{h})^{L_1'}{}_{R_1'}(\tilde{h}^*)_{L_2}{}^{R_2}(\tilde{h}^*)_{L_2'}{}^{R_2'} \quad \epsilon_{L_1 L_1'}\,\epsilon^{L_2 L_2'}\,\epsilon_{R_2 R_2'}\,\epsilon^{R_1 R_1'}
\end{aligned}
$$

$$m_h^2 = \delta v_3^2 \quad , \quad m_R^2 = \varepsilon v_3^2 \quad , \quad m_{\tilde{\Phi}}^2 = 0$$

$$\lambda_a = \lambda_1 - \lambda_3 - \varepsilon - \delta - \frac{(\lambda_1 - \lambda_3 - \varepsilon)^2}{\lambda_1 - \lambda_3 - \varepsilon - \delta} \approx -2\delta + \mathcal{O}(\varepsilon^2, \varepsilon\delta, \delta^2)$$

$$\lambda_b = \lambda_1 - \lambda_3 - \varepsilon - \delta - \frac{(\lambda_1 - \lambda_3 - \delta)^2}{\lambda_1 - \lambda_3 - \varepsilon - \delta} \approx -2\varepsilon + \mathcal{O}(\varepsilon^2, \varepsilon\delta, \delta^2)$$

$$\lambda_c = 0 \quad , \quad \lambda_d = 0 \quad , \quad \lambda_e = 0$$

$$\lambda_f = 2\left[\lambda_1 + \lambda_3 - \frac{(\lambda_1 - \lambda_3 - \varepsilon)(\lambda_1 - \lambda_3 - \delta)}{\lambda_1 - \lambda_3 - \varepsilon - \delta}\right] \approx 4\lambda_3 + \mathcal{O}(\varepsilon^2, \varepsilon\delta, \delta^2)$$

$$\lambda_g = -2\lambda_3$$

$$\lambda_h = 0$$

$$\lambda_i = \frac{1}{2}\varepsilon$$

$$\lambda_j = \frac{1}{2}\delta$$

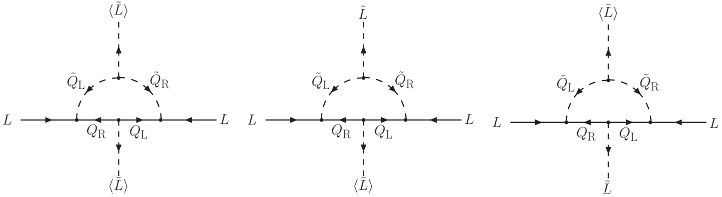
A diagrammatic procedure for matching at a higher loop order is cumbersome due to the large number of scalar fields. See arxiv:1606.07069 for tools for 1-loop matching exploiting the 1-loop effective potential.

Fermion sector in the Left-Right symmetric theory:

$$
\begin{aligned}
\mathcal{L}_{\mathrm{F}}^{(LR)} = & \; Y_\alpha \; (\tilde{R}_l^*)^R \; (l_{\mathrm{R}}^l)_R \; \Phi^s \\
& + Y_\beta \; (\tilde{R}_l^*)^R \; (l_{\mathrm{R}}^s)_R \; \Phi^l \\
& + Y_\gamma \; (\tilde{R}^I)_R \; (\mathcal{Q}_{\mathrm{R}}{}^J)^R \; D_{\mathrm{L}}^s \; \epsilon_{IJ} \\
& \\
& + Y_\delta \; (\tilde{h}^*)_L{}^R \; (H^s)^L{}_R \; \Phi^s \\
& + Y_\epsilon \; (\tilde{h}^*)_L{}^R \; (l_{\mathrm{L}}^s)^L \; (l_{\mathrm{R}}^s)_R \\
& + Y_\zeta \; (\tilde{h})^L{}_R \; (\mathcal{Q}_{\mathrm{L}}{}^I)_L \; (\mathcal{Q}_{\mathrm{R}}{}^J)^R \; \epsilon_{IJ} \\
& \\
& + Y_\eta \; \tilde{\Phi}_l^* \; \Phi^l \Phi^s \\
& \\
& + \frac{m_{\Phi^s}}{2} \Phi^s \Phi^s \qquad\qquad + \mathrm{c.c.}
\end{aligned}
$$

Tree-level matching: $Y_\gamma = -Y_\zeta = y$, and the rest vanish.

$Y_{\alpha,\beta,\epsilon}$ and $m_{\Phi^s}$ are generated at 1-loop due to the "squarks":




and $Y_\delta$ at 2-loop.
$\Rightarrow$ We get a hierarchy among the low-energy Yukawa couplings since they are generated at different loop orders!

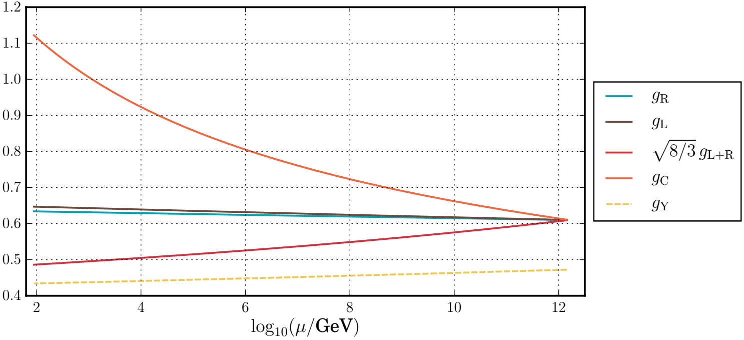


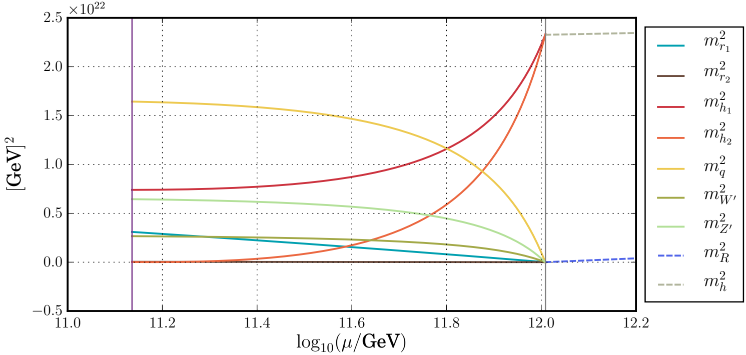

Figure: A trinification gauge coupling $g \sim 0.6$ at a matching/unification scale $\mu \sim 10^{12}\mathrm{GeV}$ leads to roughly the values of the SM gauge couplings at the electro-weak scale.

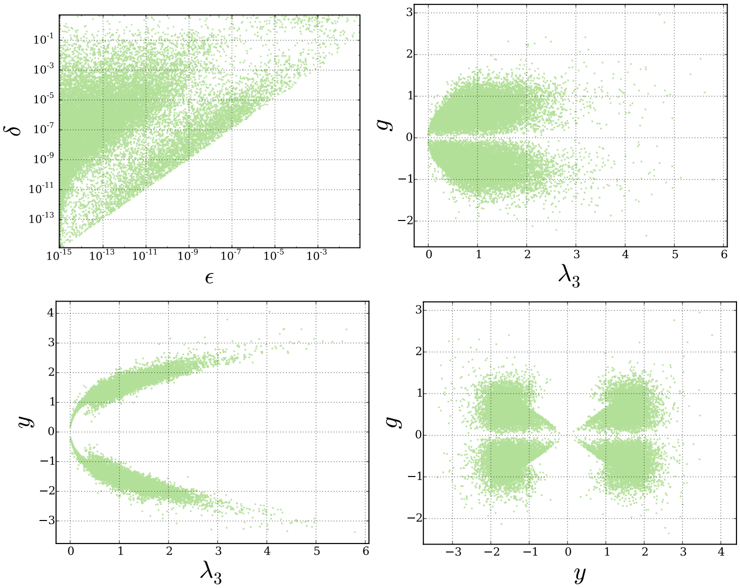
To study the model in the IR, we use the RGE to run down the model parameters. If $m_R^2$ runs negative, we generate a vev in $\tilde{R}^l{}_R$, i.e. we radiatively break $\mathrm{SU(2)_R} \times \mathrm{U(1)_{L+R}} \to \mathrm{U(1)_Y}$.

A sign flip in $m_R^2$ is possible due to the Higgs bi-doublet:

$$\begin{aligned}(4\pi)^2 \beta_{m_R^2} = {} & 4\,(2\lambda_f + \lambda_g)\, m_h^2 \\ & + \left(2|Y_\beta|^2 + 6|Y_\gamma|^2 + 2|Y_\alpha|^2 + 20\lambda_b - 8\lambda_i - 6g_{\mathrm{L+R}}^2 - \frac{9}{2}g_{\mathrm{R}}^2\right) m_R^2 \\ & - 4|Y_\alpha|^2 m_{\Phi^s}^2\end{aligned}$$

(Computation of $\beta$-functions and numerical running were performed using `pyr@te`)


$\times 10^{22}$
2.5
2.0
1.5
1.0
0.5
0.0
−0.5
$[\mathbf{GeV}]^2$
11.0
11.2
11.4
11.6
11.8
12.0
12.2
$\log_{10}(\mu/\mathbf{GeV})$
$m^2_{r_1}$
$m^2_{r_2}$
$m^2_{h_1}$
$m^2_{h_2}$
$m^2_q$
$m^2_{W'}$
$m^2_{Z'}$
$m^2_R$
$m^2_h$


δ
ϵ
10^-1
10^-3
10^-5
10^-7
10^-9
10^-11
10^-13
10^-15
g
λ3
y
λ3
g
y

**Concluding remarks:**

- Tree-level trinification breaking $\rightarrow$ The SM gauge group is generated radiatively at a much lower scale (for large regions in parameter space)
- A hierarchy in the low-scale Yukawa couplings stems from that they are generated at different loop levels
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