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Physics Letters B 637 (2006) 85–89
PHYSICS LETTERS B
www.elsevier.com/locate/physletb

3–3–1 models at electroweak scale

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Received 10 January 2006; received in revised form 26 February 2006; accepted 6 April 2006

Available online 24 April 2006

Editor: M. Cvetič

Abstract

We show that in 3–3–1 models there exist a natural relation among the $SU(3)_L$ coupling constant g, the electroweak mixing angle θ_W, the mass of the W, and one of the vacuum expectation values, which implies that those models can be realized at low energy scales and, in particular, even at the electroweak scale. So that, being that symmetries realized in Nature, new physics may be really just around the corner.

PACS: 12.10.Dm; 12.10.Kt; 14.80.Mz

Keywords: 3–3–1 model; Neutral currents

1. Introduction

Many of the extension of the Standard Model (SM) implies the existence of at least one extra neutral vector boson, say Z', which should have a mass of the order of few TeV in order to be consistent with present phenomenology. This is the case, for instance, in left–right models [1], any grand unified theories with symmetries larger than $SU(5)$ as $SO(10)$ and E_6 [2], little Higgs scenarios [3], and models with extra dimensions [4]. This makes the search for extra neutral gauge bosons one of the main goals of the next collider experiments [5]. Usually, the interactions involving Z' are parametrized (besides the pure kinetic term) as [6,7]

$$\mathcal{L}^{NC(Z')} = -\frac{\sin\xi}{2} F'_{\mu\nu} F^{\mu\nu} + M^2_{Z'} Z'_\mu Z'^\mu + \delta M^2 Z'_\mu Z^\mu - \frac{g}{2c_W} \sum_i \bar{\psi}_i \gamma^\mu \left(f^i_V - f^i_A \gamma^5\right) \psi_i Z'_\mu, \tag{1}$$

where Z, which is the would be neutral vector boson of the SM, and Z' are not yet mass eigenstates, having a mixing defined by the angle $\tan 2\phi = \delta M^2/(M^2_{Z'} - M^2_Z)$; $c_W \equiv \cos\theta_W$ (and for future use $s_W \equiv \sin\theta_W$) the usual parameter defined through the electroweak mixing angle θ_W. If Z_1 and Z_2 denote the mass eigenstates, then in most of the models $M_{Z_2} \gg M_{Z_1} \approx M_Z$. In this situation the vector and axial-vector couplings, $g_V^{(\mathrm{SM})}$ and $g_A^{(\mathrm{SM})}$, respectively, of the SM Z boson with the known fermions are modified at tree level as follows:

$$g^i_V = g^{i(\mathrm{SM})}_V c_\phi + f^i_V s_\phi, \qquad g^i_A = g^{i(\mathrm{SM})}_A c_\phi + f^i_A s_\phi, \tag{2}$$

where $g^{i(\mathrm{SM})}_V = T^i_3 - 2Q_i s^2_W$ and $g^{i(\mathrm{SM})}_A = T^i_3$, being $T^i_3 = \pm 1/2$ and Q_i the electric charge of the fermion i; we have used the notation $c_\phi(s_\phi) = \cos\phi(\sin\phi)$. The coefficients $f^i_{V,A}$ in Eq. (2) are not in general the same for all particles of the same electric charge, thus, Z' induces flavor changing neutral currents (FCNC) which imply strong constraints coming from experimental data such as ΔM_K and other $|\Delta S| = 2$ processes. These constraints imply a small value for the mixing angle ϕ or, similarly, a large value to the energy scale, generically denoted by Λ, related with the larger symmetry. If $s_\phi = 0$ is imposed such constraints could be avoided, however in most of the models with Z' this usually implies a fine tuning among $U(1)$ charges and vacuum expectation values, that is far from being natural [8].

Here we will show that there are models in which, at the tree level, it is possible that: (i) there is no mixing between Z and Z',

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0370-2693
doi:10.1016/j.physletb.2006.04.015

and the latter boson may have a mass even below the TeV scale; (ii) $\rho_0 = 1$ since $M_{Z_1} = M_Z$, and (iii) the couplings of Z_1 with fermions, $g^i_{V,A}$, being exactly those of the SM, $g^{i(\mathrm{SM})}_{V,A}$, no matter how large is Λ. This is implied not by a fine tuning but by a condition which can be verified experimentally involving the parameters of the model, g, M_W, s_W and one vacuum expectation value (VEV).

2. The model

The so-called 3–3–1 models are interesting extensions of the standard model [9–11] in which it is possible to explain the number of generations and they are also very predictive concerning new theoretical ideas as extra dimensions [12] and the little Higgs mechanism [13]. Those models also include an extra neutral vector boson so that there is, in general, a mixing of Z, the vector boson of $SU(2)_L \subset SU(3)_L$, and Z' the gauge boson related to the $SU(3)_L$ symmetry. Working in the Z, Z' basis (the parameterization in Eq. (1) is also valid but in these models there is no mixing in the kinetic term i.e., $\sin\xi = 0$) it means that the condition $\sin\phi \ll 1$ can be obtained if in this case the energy scale $\Lambda \equiv v_\chi$, with v_χ related to the $SU(3)_L$ symmetry, is above the TeV scale. Hence, it is usually believed that only approximately we can have that $Z_1 \approx Z$, even at the tree level. The same happens with the neutral current couplings, $g^i_{V,A}$, which only approximately coincide with $g^{i(\mathrm{SM})}_{V,A}$. This is true since the corrections to the Z mass and $g^i_{V,A}$ in these models, assuming $v_\chi \gg v_W \simeq 246$ GeV, are proportional to $(v_W/v_\chi)^2$ and for $v_\chi \to \infty$ we recover exactly the SM with all its degrees of freedom, with the heavier ones introduced by the $SU(3)_L$ symmetry decoupled. However, we expect that v_χ should not be extremely large if new physics is predicted to show up in the near future experiments. In practice, measurements of the ρ_0 parameter, and FCNC processes like ΔM_K, should impose constraints upon the v_χ scale at which the $SU(3)_L$ symmetry arises.

Let us consider, for instance, the model of Ref. [9] in which the electric charge operator is defined as $\mathcal{Q} = (T_3 - \sqrt{3}T_8) + X$, where T_i are the usual $SU(3)$ generators and X the charge assigned to the Abelian factor $U(1)_X$. Thus, the SM fermionic content is embedded in the extended group according to the multiplets transforming under $SU(3)_L \otimes U(1)_X$ as: for leptons, $\Psi_{aL} = (\nu_a, l^-_a, (l^-)^c_a)^T_L \sim (\mathbf{3}, 0)$, $a = e, \mu, \tau$ (the superscript c means charge conjugation operation); and for quarks, $Q_{mL} = (d_m, u_m, j_m)^T_L \sim (\mathbf{3}^*, -1/3)$; $m = 1, 2$; $Q_{3L} = (u_3, d_3, J)^T_L \sim (\mathbf{3}, 2/3)$, $u_{\alpha R} \sim (\mathbf{1}, 2/3)$, $d_{\alpha R} \sim (\mathbf{1}, -1/3)$, $\alpha = 1, 2, 3$, $j_{mR} \sim (\mathbf{1}, -4/3)$, and $J_R \sim (\mathbf{1}, 5/3)$. Here j_m and J are new quarks needed to complete the representations. To generate masses for all these fields through spontaneous symmetry breaking three triplets of Higgs scalars and a sextet are introduced; they are $\eta = (\eta^0, \eta^-_1, \eta^+_2)^T \sim (\mathbf{3}, 0)$, $\rho = (\rho^+, \rho^0, \rho^{++})^T \sim (\mathbf{3}, +1)$, $\chi = (\chi^-, \chi^{--}, \chi^0)^T \sim (\mathbf{3}, -1)$ and

$$S = \begin{pmatrix} \sigma^0_1 & h^-_1 & h^+_2 \\ h^-_1 & H^{--}_1 & \sigma^0_2 \\ h^+_2 & \sigma^0_2 & H^{++}_2 \end{pmatrix} \sim (\mathbf{6}, 0). \tag{3}$$

The VEVs in the neutral components of the scalar multiplets are defined as $\langle\eta^0_1\rangle = v_\eta/\sqrt{2}$, $\langle\rho^0_1\rangle = v_\rho/\sqrt{2}$, $\langle\chi^0_1\rangle = v_\chi/\sqrt{2}$ and $\langle\sigma^0_2\rangle = v_s/\sqrt{2}$. It is also possible to have $\langle\sigma^0_1\rangle \neq 0$ giving Majorana mass to the neutrinos, but we will not be concerned with this here. The VEV $\langle\chi^0_1\rangle$ reduces the symmetry to the SM $SU(2)_L \otimes U(1)_Y$ symmetry and the other VEVs further reduce it to the electromagnetic $U(1)_Q$ factor.

From the kinetic terms for the scalar fields, constructed with the covariant derivatives

$$\mathcal{D}_\mu\varphi = \partial_\mu\varphi - ig\vec{W}_\mu\cdot\vec{T}\varphi - ig_X X\varphi B_\mu,$$
$$\mathcal{D}_\mu S = \partial_\mu S - ig\big[\vec{W}_\mu\cdot\vec{T}S + S\vec{W}_\mu\cdot\vec{T}^T\big], \tag{4}$$

where g_X denotes the $U(1)_X$ gauge coupling constant and $\varphi = \eta, \rho, \chi$, we obtain the mass matrices for the vector bosons. Besides $W^\pm$ there are two other charged vector bosons, $V^\pm$ and $U^{\pm\pm}$. The masses of these charged vector bosons are given exactly by $M^2_W = (g^2/4)v^2_W$, $M^2_V = (g^2/4)(v^2_\eta + 2v^2_s + v^2_\chi)$ and $M^2_U = (g^2/4)(v^2_\rho + 2v^2_s + v^2_\chi)$, where we have defined $v^2_W = v^2_\eta + v^2_\rho + 2v^2_s$ (in models where there are heavy leptons transforming nontrivially under $SU(3)_L \otimes U(1)_X$ there is not the contribution of the sextet and the above equations are still valid simply doing $v_s = 0$ [10]). For the mass square matrix of the neutral vector bosons in this model we have the following form, after defining the dimensionless ratios $\bar{v}_\rho = v_\rho/v_\chi$, $\bar{v}_W = v_W/v_\chi$ and the parameter $t^2 = g^2_X/g^2 = s^2_W/(1 - 4s^2_W)$,

$$\mathcal{M}^2 = \frac{g^2}{4}v^2_\chi \begin{pmatrix} \bar{v}^2_W & \frac{1}{\sqrt{3}}(\bar{v}^2_W - 2\bar{v}^2_\rho) & -2t\bar{v}^2_\rho \\ \frac{1}{\sqrt{3}}(\bar{v}^2_W - 2\bar{v}^2_\rho) & \frac{1}{3}(\bar{v}^2_W + 4) & \frac{2}{\sqrt{3}}t(\bar{v}^2_\rho + 2) \\ -2t\bar{v}^2_\rho & \frac{2}{\sqrt{3}}t(\bar{v}^2_\rho + 2) & 4t^2(\bar{v}^2_\rho + 1) \end{pmatrix}, \tag{5}$$

in the $(W^3_\mu, W^8_\mu, B_\mu)$ basis. This matrix has a zero eigenvalue corresponding to the photon and two nonzero ones which are given by

$$M^2_{Z_1} = \frac{g^2 v^2_\chi}{6}\big[3t^2\big(\bar{v}^2_\rho + 1\big) + 1 + \bar{v}^2_W\big](1 - R), \tag{6}$$

$$M^2_{Z_2} = \frac{g^2 v^2_\chi}{6}\big[3t^2\big(\bar{v}^2_\rho + 1\big) + 1 + \bar{v}^2_W\big](1 + R), \tag{7}$$

with

$$R = \left[1 - \frac{3(4t^2 + 1)(\bar{v}^2_W(\bar{v}^2_\rho + 1) - \bar{v}^4_\rho)}{(3t^2(\bar{v}^2_\rho + 1) + 1 + \bar{v}^2_W)^2}\right]^{1/2}. \tag{8}$$

3. ρ_1 and ρ_0 parameters

In order to analyze the condition which allows to identify Z_1 of the 3–3–1 model with the Z of the SM, let us introduce a dimensionless ρ_1-parameter defined at the tree level as $\rho_1 = c^2_W M^2_{Z_1}/M^2_W$. As we can see from Eqs. (6) and (7) both mass eigenvalues, M_{Z_1} and M_{Z_2}, have a complicate dependence on the VEVs but we observe from Eq. (6) that $\rho_1 \leqslant 1$ (or, $M_{Z_1} \leqslant M_Z$) is a prediction of the model. Next, we can search for the conditions under which we have $\rho_1 \equiv \rho_0 = 1$, where

$\rho_0 = c_W^2 M_Z^2/M_W^2$ is the respective parameter in the SM. This is equivalent to the condition that $M_{Z_1} \equiv M_Z$ at the tree level. The equation $\rho_1 = 1$ has besides the solution $v_\chi \to \infty$, another less trivial one which can be obtained using Eq. (6) above:

$$\bar{v}_\rho^2 = \frac{1-4s_W^2}{2c_W^2}\bar{v}_W^2. \tag{9}$$

The condition in Eq. (9) implies, using the definition of v_W given above also $v_\eta^2 + 2v_s^2 = [(1+2s_W^2)/2c_W^2]v_W^2$. We recall that the $U(1)_X$ quantum number of the η and S fields are different from that of the ρ field, so there is no symmetry among the respective VEVs. We have verified that (9) is stable in the following sense: small deviations from it implies small deviations from $\rho_1 = 1$. With $s_W^2 = 0.2312$ [7] we obtain $v_\rho \approx 54$ GeV and $\sqrt{v_\eta^2 + 2v_s^2} \approx 240$ GeV. Notice that Eq. (9) is independent of the v_χ scale. Hence, all consequences of it will be also independent of v_χ as claimed above in the Introduction. The fact that v_χ does not need to have a large value to be consistent with the present phenomenology is interesting in the models of Refs. [9,10] since these models have a Landau-like pole at the TeV scale [14].

If we substitute Eq. (9) in Eqs. (6) and (7) we obtain $M_{Z_1}^2 = (g^2/4c_W^2)v_W^2 \equiv M_Z$ and

$$M_{Z_2}^2 \equiv M_{Z'}^2 = \frac{g^2 v_W^2}{2}\frac{(1-2s_W^2)(4+\bar{v}_W^2) + s_W^4(4-\bar{v}_W^4)}{6c_W^2(1-4s_W^2)}v_\chi^2. \tag{10}$$

Thus, assuming that Eq. (9) is valid we shall not distinguish between Z_1 and Z and between Z_2 and Z' unless stated explicitly. Moreover, the mass of Z' can be large even if v_χ is of the order of the electroweak scale. In fact, from Eq. (10) we see that for $\bar{v}_W = 1$ (the electroweak scale is equal to the 3–3–1 scale) we obtain $M_{Z'} = 3.77 M_W$. Of course for lower values of $\bar{v}_W$, Z' is heavier, for instance for $\bar{v}_W = 0.25$ we have $M_{Z'} = 18.36 M_W$. We recall that since v_χ does not contribute to the W mass it is not constrained by the 246 GeV upper bound. Thus, independently if $\bar{v}_W^2$ is larger, smaller or equal to 1, the charged vector boson V is heavier than U, being $\Delta M = \sqrt{M_V^2 - M_U^2} = 75.96$ GeV, when $v_s = 0$.

4. Neutral current couplings

We have also obtained the full analytical exact expressions for the neutral current couplings $g_{V,A}^i$ and $f_{V,A}^i$, and verified that they also depend on the VEVs in a complicated way. But when Eq. (9) is used in those expressions we obtain for the case of the known fermions $g_{V,A}^i \equiv g_{V,A}^{i(\mathrm{SM})}$, and $f_{V,A}^i = f_{V,A}^i(s_W)$, i.e., these couplings depend only on the electroweak mixing angle. For the lepton couplings with Z', also after using Eq. (9) in the general expressions, we obtain $f_V^\nu = f_A^\nu = f_V^l = -f_A^l = -\sqrt{3(1-4s_W^2)}/6 (\approx -0.07)$. We see that the couplings for all leptons with Z' are leptophobic [15]. In particular, the couplings of Z to the exotic quarks j_m and J are given by $g_V^{j_m} = (8/3)s_W^2$, $g_V^J = -(10/3)s_W^2$ and $g_A^{j_m} = g_A^J = 0$. Notice also that the exotic quarks have pure vectorial couplings with Z. The couplings of Z' in the quark sector are given by:

$$f_V^{u_m} = \frac{1}{2\sqrt{3}}\frac{1-6s_W^2}{\sqrt{1-4s_W^2}}, \qquad f_A^{u_m} = \frac{1}{2\sqrt{3}}\frac{1+2s_W^2}{\sqrt{1-4s_W^2}},$$
$$f_V^{u_3} = -\frac{1}{2\sqrt{3}}\frac{1+4s_W^2}{\sqrt{1-4s_W^2}}, \qquad f_A^{u_3} = -\frac{1}{\sqrt{3}}\sqrt{1-4s_W^2},$$
$$f_V^{d_m} = \frac{1}{2\sqrt{3}\sqrt{1-4s_W^2}}, \qquad f_A^{d_m} = \frac{\sqrt{1-4s_W^2}}{2\sqrt{3}},$$
$$f_V^{d_3} = -\frac{1}{2\sqrt{3}}\frac{1-2s_W^2}{\sqrt{1-4s_W^2}}, \qquad f_A^{d_3} = -\frac{1}{2\sqrt{3}}\frac{1+2s_W^2}{\sqrt{1-4s_W^2}},$$
$$f_V^{j_m} = -\frac{1}{\sqrt{3}}\frac{1-9s_W^2}{\sqrt{1-4s_W^2}}, \qquad f_A^{j_m} = -\frac{1}{\sqrt{3}}\frac{c_W^2}{\sqrt{1-4s_W^2}},$$
$$f_V^J = \frac{1}{\sqrt{3}}\frac{1-11s_W^2}{\sqrt{1-4s_W^2}}, \qquad f_A^J = \frac{1}{\sqrt{3}}\frac{c_W^2}{\sqrt{1-4s_W^2}}. \tag{11}$$

In literature [9,16] these couplings were considered as an approximation of the exact couplings. Notice that, all these couplings refer to fermions which are still symmetry eigenstates, thus we see that in the leptonic sector there are not FCNCs neither with Z nor with Z' and, in the quark sector there are FCNC only coupled to Z' as can be see from Eq. (11). However, FCNC mediated by the Z' depend only on its mass, but these FCNC are not necessary large since there are also contributions in the scalar sector (see below).

The main feature introduced by the validity of the condition Eq. (9), which we would like to stress, is the fact that all couplings between the already known particles are exactly those of the SM, regardless the value of the v_χ scale. Hence, v_χ is not required to be large to recover those observed couplings (until now $v_\chi \to \infty$ was the usual approach to do that). In this way, the 3–3–1 gauge symmetry could be realized, for instance, at the electroweak scale ($v_\chi = v_W$) allowing the extra particles introduced by the $SU(3)_L$ to be light enough to not decouple and be discovered in the near future experiments.

5. A Goldberger–Treiman-like relation

We can rewrite Eq. (9) as

$$g\frac{v_\rho}{\sqrt{2}} = \frac{\sqrt{1-4s_W^2}}{c_W}M_W. \tag{12}$$

This is like the Goldberger–Treiman relation [17] in the sense that its validity implies a larger symmetry of the model (see below) and all quantities appearing in it can be measured independently of each other. In fact, all but v_ρ, are already well known. However, cross sections of several processes, for instance $e^+e^- \to ZH$ where H is a neutral Higgs scalar transforming as doublet of $SU(2)$, are sensitive to the value of v_η (or v_ρ) [18]. So, in principle it is possible to verified if Eq. (9), or equivalently Eq. (12), is satisfied and if the 3–3–1 symmetry can be implemented near the weak scale.

6. Custodial symmetries and the oblique T parameter

We can understand the physical meaning of Eq. (9) in the following way. The 3–3–1 models have an approximate $SU(2)$ custodial symmetry. This is broken by the mixing between Z and Z'. In general we have a mixture between these neutral bosons in such a way that the mass eigenstates Z_1 and Z_2 can be written as [19] $Z_1 = Zc_\phi - Z's_\phi$ and $Z_2 = Zs_\phi + Z'c_\phi$, and the condition in Eq. (9) is equivalent to put $\phi = 0$ i.e., no mixing at all between Z and Z'. There is also an approximate $SU(3)$ custodial symmetry because when both (9) and $\sin\theta_W = 0$ are used, we have $M_U/M_{Z'} = 1$. However this symmetry is badly broken. We stress that the alternative approach used in literature [9,19,20] is that the condition $\sin\phi \ll 1$ is obtained by assuming that $v_\chi \gg v_W$. This is of course still a possibility if the relation (9) is not confirmed experimentally. However, we have shown above that it is possible that $\phi = 0$ even if $v_\chi = v_W$.

Of course, in any case radiative corrections will induce a mixing among Z and Z', i.e., a finite contribution to ϕ. This should imply small deviations from $\rho_0 = 1$. The oblique T parameter constraints this deviations since $\rho_0 - 1 \simeq \alpha T$ and it is given, for the 3–3–1 models, in Ref. [21]. Using the expressions of Ref. [21] but without the mixing at the tree level ($\phi = 0$ in Eq. (4.1) of [21]), we obtain for example $T = -0.1225$ for $\bar{v}_W = 1$ and $T = -0.012$ for $\bar{v}_W = 0.25$, with $T \to 0$ as $\bar{v}_W \to 0$ ($v_\chi \to \infty$), and all T values calculated with $\bar{v}_W \leqslant 1$ are within the allowed interval [7]. This implies that the condition Eq. (9) is not significantly disturbed by radiative corrections. It means that the natural value of $\sin\phi$, arisen only through radiative corrections, is small because the symmetry of the model is augmented when this parameter vanishes.

The important thing is that even if Eq. (9) or equivalently Eq. (12) are valid only approximately, we will have that again $M_{Z_1} \approx M_Z$ and also the neutral current couplings of Z_1 only approximately coincide with those of the SM, but now this is valid almost independently of the value of v_χ. That is, the 3–3–1 symmetry still can be implemented at an energy scale near the electroweak scale.

7. Experimental constraints on the $SU(3)_L$ scale

Once the v_χ scale is arbitrary when Eq. (9) is satisfied, we can ask ourselves what about the experimental limit upon the masses of the extra particles that appear in the model. After all they depend mainly on v_χ, the scale at which the $SU(3)_L$ symmetry is supposed to be valid. Firstly, let us consider the Z' vector boson. It contributes to the ΔM_K at the tree level [22]. This parameter imposes constraints over the quantity $(\mathcal{O}^d_L)_{3d}(\mathcal{O}^d_L)_{3s}(M_Z/M_{Z'})$, which must be of the order of 10^{-4} to have compatibility with the measured ΔM_K. This can be achieved with $M_{Z'} \sim 4$ TeV if we assume that the mixing matrix have a Fritzsch-structure $\mathcal{O}^d_{Lij} = \sqrt{m_j/m_i}$ [23] or, it is possible that the product of the mixing angles saturates the value 10^{-4} [22], in this case Z' can have a mass near the electroweak scale. More important is the fact that there are also in this model FCNC mediated by the neutral Higgs scalar which contributes to ΔM_K. These contributions depend on the mixing matrix in the right-hand d-quark sector, $\mathcal{O}^d_R$, and also on some Yukawa couplings, Γ^d, i.e., the interactions are of the form $\bar{d}_L(\mathcal{O}^d_L)_{d3}\Gamma^d_{3\alpha}(\mathcal{O}_R)_\alpha s_R$. Thus, their contributions to ΔM_K may have opposite sing relative to that of the contribution of Z'. A realistic calculation of the ΔM_K in the context of 3–3–1 models has to take into account these extra contributions as well. Muonium–antimuonium transitions would imply a lower bound of 850 GeV on the masses of the doubly charged gauge bileptons, U^{--} [24]. However this bound depends on assumptions on the mixing matrix in the lepton charged currents coupled to U^{--} and also it does not take into account that there are in the model doubly charged scalar bileptons which also contribute to that transition [25]. Concerning these doubly charged scalars, the lower limit for their masses are only of the order of 100 GeV [26]. From fermion pair production at LEP and lepton flavor violating effects suggest a lower bound of 750 GeV for the mass M_U but again it depends on assumptions on the mixing matrix and on the assumption that those processes are induced only by the U^{--} boson [27]. Other phenomenological analysis in e^+e^-, $e\gamma$ and $\gamma\gamma$ colliders assume bileptons with masses between 500 GeV and 1 TeV [28,29]. The muonium fine structure only implies $M_U/g > 215$ GeV [30] but also ignores the contributions of the doubly charged scalars. Concerning the exotic quark masses there is no lower limit for them but if they are in the range of 200–600 GeV they may be discovered at the LHC [31]. Direct search for quarks with $Q = (4/3)$e imply that they are excluded if their mass is in the interval 50–140 GeV [32] but only if they are stable. Similarly, most of the searches for extra neutral gauge bosons are based on models that do not have the couplings with the known leptons and quarks as those of the 3–3–1 model [6], anyway we have seen that even if $v_\chi = v_W$ the Z' has a mass of the order of 300 GeV. Finally, rare processes like $\mu^- \to e^-\nu_e\nu^c_\mu$, induced by the extra particles are also not much restrictive. We may conclude that there are not yet definitive experimental bounds on the masses of the extra degrees of freedom of the 3–3–1 models.

8. Conclusions

Summarizing, we have shown that if the condition in Eq. (9) is realized, concerning the already known particles, the 3–3–1 model and the SM are indistinguishable from each other at the tree level and, as suggested by the value of the T-parameter obtained above, it is possible that this happens even at the one loop level. The models can be confirmed or ruled out by searching directly for the effects of their new particles, for instance in left–right asymmetries in lepton–lepton scattering. An asymmetry of this sort only recently has begun to be measured in an electron–electron fixed target experiment [33], but the effects of these asymmetries could be more evident in collider experiments [34,35]. From all we have discussed above, it is clear that new physics may be really just around the corner.

We have also verified that in 3–3–1 models with heavy leptons [10] and with right-handed neutrinos transforming nontrivially under the 3–3–1 gauge symmetry [11] a similar situation occurs, but, in the later model, the equivalent of the relation in Eq. (9) is given by $\bar{v}_\rho^2 = [(1-2s_W^2)/2c_W^2]\bar{v}_W^2$ [36]. Notwithstanding we have verified that the latter condition is less stable, in the sense we said before: small deviation of it implies large deviation from $\rho_1 = 1$. This suggests that, when both 3–3–1 models were embedding in a $SU(4)_L \otimes U(1)_N$ model [37], the $SU(3)$ subgroup which contains the vector boson Z of the SM should be the minimal 3–3–1 model considered in this work.

Acknowledgements

A.G.D. is supported by FAPESP under the process 05/52006-6, and this work was partially supported by CNPq under the processes 305185/03-9 (J.C.M.), and 306087/88-0 (V.P.).

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