Vol. 23(1992) ACTA PHYSICA POLONICA B No 10

THE CONTINUOUS LEFT-RIGHT SYMMETRIC AND ANTISYMMETRIC MODELS*

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(Received September 28, 1992)

The continuous left-right (LR) model which is characterized by the orientation angle of $SU(2)_R$ generator in the group space is suggested. This model reproduces all the known LR models. The cross sections of the processes

$$Z_n \to f\bar{f}\,,\ pp \to Z_n + X\,,\ p\bar{p} \to Z_n + X\,,\ f\bar{f} \to W_k^- W_l^+ \qquad (n,k,l = 1,2)$$

are obtained. The comparison of the symmetric version of this model with the experiments is fulfilled.

PACS numbers: 12.15. -y

1. Introduction

The parity violation (PV) in weak interactions is one of the puzzles of modern physics which cannot get the satisfactory explanation within the standard model of electroweak interactions (SM) in which the gauge symmetry is broken spontaneously while the parity is broken explicitly. The most attractive enlargement of the SM is the left-right symmetric (LR) model based on the $SU(2)_L \times SU(2)_R \times U(1)$ gauge group, where the observed PV in weak interactions is connected with the mechanism of the spontaneous symmetry breaking. The PV could be caused by the difference in the $SU(2)_L$ and $SU(2)_R$ sectors between the following quantities: (a) elements of Kobayashi–Maskawa matrix; (b) couplings constants $g(SU(2)_L) \equiv g_L$ and $g(SU(2)_R) \equiv g_R$; (c) gauge boson masses; (d) Majorana neutrinos

* Presented at the XXXII Cracow School of Theoretical Physics, Zakopane, Poland, June 2–12, 1992.

masses. The LR model also has some other attractive features: (a) quantum numbers of group U(1) are identified with B–L what allows to connect the violation of the parity and local B–L symmetry (where B and L are baryon and lepton numbers); (b) the theory due to see-saw mechanism provide the existence of light left-handed and heavy right-handed neutrinos; (c) the effect of the weak CP violation both in quark and lepton sectors is caused by the mechanism of spontaneous symmetry breaking and its values (the ratio of decay amplitudes of K_l^0 and K_s^0 mesons into pions, a charge asymmetry parameter of lepton decays of K_l^0 mesons, electric dipole moments of a neutron and charged leptons, *etc.*) are the same order as it follows from the experiment.

In this work both the symmetric and asymmetric left-right (LR) theories are investigated. There are a lot of papers in which such models are considered [1–5]. Many of them are discriminated one from the other by the choice of the transformation to mass eigenstate basis in the potentials space of the neutral gauge bosons. This transformation is determined by the Higgs sector structure and can be written as the product of two matrix Λ and U_N which carry out the transition from the initial basis to the final one by the chain

$$\begin{pmatrix} W_{3\mu}^{L} \\ W_{3\mu}^{R} \\ B_\mu \end{pmatrix} \xrightarrow{\Lambda} \begin{pmatrix} Z_\mu^{L} \\ Z_\mu^{R} \\ A_\mu \end{pmatrix} \xrightarrow{U_N} \begin{pmatrix} Z_{1\mu} \\ Z_{z\mu} \\ A_\mu \end{pmatrix},$$

where $W_{3\mu}^{L,R}$ and B_μ are the gauge fields corresponding to the SU(2) and U(1) groups, respectively. The U_N has the same form in all the theories

$$U_N = \begin{pmatrix} c_\Phi & s_\Phi & 0 \\ -s_\Phi & c_\Phi & 0 \\ 0 & 0 & 1 \end{pmatrix}$$

while the expression for the matrix Λ might not coincide in the different versions of the LR models. For example, two versions of the symmetric LR models proposed in Refs [1, 2] (further I shall call them LR1 and LR2) have the matrix Λ which is defined by the following expressions

$$\Lambda_1 = \begin{pmatrix} \frac{1}{\sqrt{2}} & \frac{-1}{\sqrt{2}} & 0 \\ \sqrt{\frac{(c_W^2 - s_W^2)}{2}} & \sqrt{\frac{(c_W^2 - s_W^2)}{2}} & \frac{-s_W\sqrt{2}}{} \\ s_W & s_W & \sqrt{c_W^2 - s_W^2} \end{pmatrix}$$

in LR1 and

$$\Lambda_2 = \begin{pmatrix} c_W & \frac{-s_W^2}{c_W} & \frac{-s_W\sqrt{c_W^2 - s_W^2}}{c_W} \\ 0 & \sqrt{\frac{c_W^2 - s_W^2}{c_W}} & \frac{-s_W}{c_W} \\ s_W & s_W & \sqrt{c_W^2 - s_W^2} \end{pmatrix}$$

LR2 respectively, where $c_\Phi = \cos\phi$, $s_\Phi = \sin\Phi$, $c_W = \cos\theta_W$, $s_W = \sin\theta_W$, Φ is the mixing angle of the neutral gauge bosons. For our analysis it is convenient to unify all both symmetric and asymmetric models in the one common model. That could be achieved by the corresponding parametrization of the matrix Λ. In Section 2 a theoretical description of such a model is provided. In Section 3 we analyse the model from the point of view of experiments coming from LEP I and future colliders. Finally, the work is summarized in Section 4.

2. Description of the model

Now, I should like to make some remarks on the electroweak models. In the SM the choice of the gauge group fixes the form of both the neutral current Lagrangian $\mathcal{L}_{\rm NC}$ and the Lagrangian describing the trilinear-gauge-boson couplings (TBC) $\mathcal{L}_{WWZ}$. In superstring motivated extra U(1) models the situation is not the same. In the most general case there are two extra Z-bosons called Z_ψ and Z_χ. The former arises when E_6 breaks down to the SO(10)×U(1)$_\chi$ and the latter arises when the SO(10) breaks down to the SU(5)×U(1)$_\psi$. The mass eigenstates of the new gauge bosons are

$$\begin{aligned} Z_z &= Z_\psi \cos\beta + Z_\chi \sin\beta \\ Z_n &= -Z_\psi \sin\beta + Z_\chi \cos\beta\,, \end{aligned} \tag{2.1}$$

where the angle β being dependent on the vacuum expectation values (VEV) of the Higgs fields and the gauge couplings $g_{\psi,\chi}$ determines the orientation of the U'(1) generator in E_6 space. The choice of the β fixes both of $\mathcal{L}_{\rm NC}$ and $\mathcal{L}_{WWZ}$. Now depending upon the chosen value β we have the following models: (a) Ψ-model, ($\beta = 0°$); (b) X-model, ($\beta = -90°$); (c) η-model, ($\beta = 37.36°$); (d) I-model, ($\beta = 52.24°$). In literature it also considered the continuous extra U(1) model (β arbitrary) [6]. It should be noted that the parameter β redistributes the roles between the mixing angle Φ and the masses of new neutral bosons m_{Z_n} ($n = 2, 3$). As is well known it is displayed in distinction of the experimental constraints on Φ and m_{Z_2} for different extra U(1) models. It is natural to wait for arbitrariness of the generators orientation in the group space and accordingly the appearance of the analogous parameter in any electroweak theory with an extended gauge group. We face the situation when E_6 breaks down to the $SU(2)_L \times SU(2)_R \times U(1)_L \times U(1)'_R$. In above mentioned examples the kind of the transformation to the mass eigenstate basis of the neutral gauge bosons is fixed by the structure of the grand unified theory (GUT) to be exact with the earlier stage of the symmetry breaking. Up to now we can not give preference to the definite GUT with confidence. Then, investigating the $SU(2)_L \times SU(2)_R \times U(1)$ gauge theory as a low energy approximation of GUT we should consider the most general parametrization of the matrix Λ. The matrix Λ is the product of three-single parameter orthogonal transformations in the space of the potentials $W^L_{3\mu}$, $W^R_{3\mu}$, B_μ. Therefore, it has the form

$$\Lambda = \begin{pmatrix} c_{\varphi_3}c_{\varphi_1} - s_{\varphi_3}s_{\varphi_1}s_{\varphi_2} & c_{\varphi_3}s_{\varphi_1} + s_{\varphi_3}c_{\varphi_1}s_{\varphi_2} & s_{\varphi_3}c_{\varphi_2} \\ -c_{\varphi_2}s_{\varphi_1} & c_{\varphi_2}c_{\varphi_1} & -s_{\varphi_2} \\ -c_{\varphi_1}s_{\varphi_3} - c_{\varphi_3}s_{\varphi_1}s_{\varphi_2} & -s_{\varphi_1}s_{\varphi_3} + c_{\varphi_3}c_{\varphi_1}s_{\varphi_2} & c_{\varphi_3}c_{\varphi_2} \end{pmatrix}, \tag{2.2}$$

where $c_{\varphi_i} = \cos\varphi_i$, $s_{\varphi_i} = \sin\varphi_i$, the angles φ_i are the functions of the VEV's and the gauge constants.

The relation

$$Q = T_{3L} + T_{3R} + \frac{B-L}{2} \tag{2.3}$$

ensuring the conservation of $U(1)_{em}$ symmetry and masslessness of a photon fixes the third line of the matrix Λ. It gives the set of equations

$$\begin{aligned} -c_{\varphi_1}s_{\varphi_3} - c_{\varphi_3}s_{\varphi_1}s_{\varphi_2} &= \frac{e}{g_L} \\ -s_{\varphi_1}s_{\varphi_3} + c_{\varphi_3}c_{\varphi_1}s_{\varphi_2} &= \frac{e}{g_R} \\ c_{\varphi_3}c_{\varphi_2} &= \frac{e}{g'}, \end{aligned} \tag{2.4}$$

where $s_W = \sin\theta_W$, $c_W = \cos\theta_W$ and the electric charge is defined by the equation

$$e = \frac{g_L g_R g'}{\sqrt{g_L^2 g_R^2 + g'(g_L^2 + g_R^2)}}. \tag{2.5}$$

It is easy to see that only two equations in (2.4) are linear-independent. Then we can exclude φ_2 and φ_3, for example, what gives the following expression for the matrix Λ

$$\Lambda = \begin{pmatrix} e(b^{-1}c_\varphi + ba_+a_-s_\varphi) & e(b^{-1}s_\varphi - ba_+a_-c_\varphi) & -ebg'^{-1}a_+ \\ -bg'^{-1}s_\varphi & bg'^{-1}c_\varphi & -ba_- \\ \frac{e}{g_L} & \frac{e}{g_R} & \frac{e}{g'} \end{pmatrix}, \tag{2.6}$$

where $a_+ = g_R^{-1}s_\varphi + g_L^{-1}c_\varphi$, $a_- = g_R^{-1}c_\varphi - g_L^{-1}s_\varphi$, $b = \sqrt{g'^{-2} + a_-^2}$ and we suppress the index 3 at φ. So, we have the model with parameter φ.

Changing φ we can reproduce all the known LR models. Further on, we shall name this model the continuous symmetric LR model (CSLR) when $g_{\rm L} = g_{\rm R}$ or the continuous asymmetric one (CALR) when $g_{\rm L} \neq g_{\rm R}$. From (2.6) it follows that the LR1 is reproduced at $\varphi = -\pi/4$, and LR2 does at $\varphi = 0$.

Let us consider the situation when we begin to build the electroweak theory from the $SU(2)_{\rm L}\times SU(2)_{\rm R}\times U(1)$ group directly. The Lagrangian of the theory has the form

$$\begin{aligned}\mathcal{L} = &-\tfrac{1}{4}\left(W^a_{{\rm L}\mu\nu}W^a_{{\rm L}\mu\nu} + W^a_{{\rm R}\mu\nu}W^a_{{\rm R}\mu\nu} + B_{\mu\nu}B_{\mu\nu}\right)\\ &+ i\sum_f\left(\bar{\Psi}^{(f)}_{\rm L}D_\mu\gamma_\mu\Psi^{(f)}_{\rm L} + \bar{\Psi}^{(f)}_{\rm R}D_\mu\gamma_\mu\Psi^{(f)}_{\rm R}\right) + \sum_i|D_\mu\varphi_i|^2 + \mathcal{L}_Y - V\,,\end{aligned} \tag{2.7}$$

where D_μ are the usual covariant derivatives $W^a_{{\rm L}\mu\nu}$, $W^a_{{\rm R}\mu\nu}$, $B_{\mu\nu}$ denote the $SU(s)_{\rm L}$, $SU(2)_{\rm R}$ and $U(1)$ gauge fields, $\mathcal{L}_Y$ is the Lagrangian describing the interaction between Higgs particles φ_i and fermions, V is the Higgs potential. In the most general case the V can be represented in the form

$$V = \sum_i \lambda_i V^{(2)}(\varphi_i,\varphi_i) + \sum_{i,j,k,l}\mu_{ijkl}V^{(4)}(\varphi_i,\varphi_j,\varphi_k,\varphi_l)\,, \tag{2.8}$$

where λ_i and μ_{ijkl} are constants, $V^{(2)}(\varphi_i,\varphi_i)$ and $V^{(4)}(\varphi_i,\varphi_j,\varphi_k,\varphi_l)$ are the quadratic and the bi-quadratic on φ_i terms (their obvious form is defined by the symmetries imposed on (2.8)). We introduce an arbitrary number of the following Higgs multiplets (in brackets their numbers $T_{\rm L}$, $T_{\rm R}$, $(B-L)/2$ are given): (a) doublets $X_{\rm L}(\frac{1}{2},0,-\frac{1}{2})$, $X_{\rm R}(0,\frac{1}{2},-\frac{1}{2})$; (b) triplets $\delta_{\rm L}(1,0,1)$, $\delta_{\rm R}(0,1,1)$; (c) bi-doublets $\Phi(\frac{1}{2},\frac{1}{2},0)$. A minimum of the potential V corresponds to the following choice of VEV's of Higgs fields

$$\langle X_{\rm L,R}\rangle = \begin{pmatrix}0\\ v_{\rm L,R}\end{pmatrix},\quad \langle\delta_{\rm LR}\rangle = \begin{pmatrix}0\\0\\ \Delta_{\rm LR}\end{pmatrix},\quad \langle\Phi\rangle = \begin{pmatrix}k & 0\\ 0 & k'\end{pmatrix}. \tag{2.9}$$

The neutral-boson $(\text{mass})^2$ matrix is

$$M_N = \begin{pmatrix}A_{\rm L} & B & C_{\rm L}\\ B & A_{\rm R} & C_{\rm R}\\ C_{\rm L} & C_{\rm R} & G\end{pmatrix}, \tag{2.10}$$

where

$$\begin{gathered}A_{\rm LR} = g^2_{\rm L,R}(\beta_{\rm L,R}+\alpha)\,,\quad C_{\rm L,R} = -gg'_{\rm L,R}\beta_{\rm LR}\,,\quad B = -g_{\rm L}g_{\rm R}\alpha\,,\\ G = g'^2(\beta_{\rm L}+\beta_{\rm R})\,,\quad \alpha = \tfrac{1}{2}\sum\left(|k|^2+|k'|^2\right),\\ \beta_{\rm L,R} = \tfrac{1}{2}\sum|v_{\rm L,R}|^2 + 2\sum|\Delta_{\rm L,R}|^2\,.\end{gathered}$$

The matrix Λ which together with U_N diagonalizes M_N according to

$$U_N \Lambda M_N (U_N \Lambda)^T$$

is defined by the expression (2.6) both for $g_L \neq g_R$ and for $g_L = g_R$. In the former case all the quantities m_{z_1}, m_{z_2}, Φ, φ, θ_W parametrizing the neutral current interaction

$$\alpha,\ \beta_L,\ \beta_R,\ g_L,\ g_R,\ g' \longrightarrow m_{z_1},\ m_{z_2},\ \theta_W,\ \Phi,\ \varphi$$

are independent. In the latter one the φ is the function on m_{z_1}, m_{z_2}, Φ, θ_W and its values lie in the interval from 0 to $-\pi/4$.

Now one might determine the interaction Lagrangian in this model. After straightforward (though somewhat tedious) calculations we obtain the following expressions

$$\mathcal{L}_{wwv} = i\Big(W^*_{k\mu\nu}\rho^{(V)}_{kl}V_\nu + W^*_{k\nu}\rho^{(V)}_{kl}V_{\nu\mu}\Big)W_{l\mu}\,, \tag{2.11}$$

$$\mathcal{L}_{NC} = i\sum_f \bar{\Psi}_f\gamma_\mu\Big[QA_\mu + \tfrac{1}{2}\sum_{n=1}^{2}\big(g^f_{Vn} + g^f_{An}\gamma_5\big)Z_{n\mu}\Big]\Psi_f\,, \tag{2.12}$$

$$\mathcal{L}_{CC} = \frac{ig_L}{2\sqrt{2}}\sum_{\substack{j=e,d\\ i=\nu,\mu}}\left[K_{ij}\bar{\Psi}_j\gamma_\mu(1+\gamma_5)\Psi_i W^L_\mu + (L\leftrightarrow R)\right], \tag{2.13}$$

where

$$V_{\mu\nu} = \partial_\mu V_\nu - \partial_\nu V_\mu \quad (V_\mu = A_\mu,\ Z_{n\mu}) \quad n = 1,2\,,$$

$$\rho^{(A)}_{kl} = Q\delta_{kl}\,, \tag{2.14}$$

$$\rho^{(Z_1)}_{ll} = \cos^2(\xi + \tfrac{\pi}{2}\delta_{l2})g_L D_{11} + \sin^2(\xi + \tfrac{\pi}{2}\delta_{l2})g_R D_{12}\,, \tag{2.15}$$

$$\rho^{(Z_1)}_{kl} = \rho^{(Z_1)}_{lk} = \tfrac{1}{2}\sin 2\xi(g_L D_{11} - g_R D_{12}) \quad k \neq l\,, \tag{2.16}$$

$$g^f_{Vn} = \big[T^f_{3L}g_L D_{n1} + T^f_{3R}g_R D_{n2} + \big(1 - \tfrac{\delta_{f\nu}}{2}\big)(B-L)g' D_{n3}\big]\,, \tag{2.17}$$

$$g^f_{An} = \big[T^f_{3L}g_L D_{n1} - T^f_{3R}g_R D_{n2} - \delta_{f\nu}g'(T^f_{3L} - T^f_{3R})D_{n3}\big]\,. \tag{2.18}$$

D_{nm} and K_{ij} are the elements of the matrix $U_N\Lambda$ and the Kobayashi–Maskawa one, respectively, Q is the charge of the particle, ξ is the mixing angle of charged gauge bosons, and the expressions for $\rho^{(Z_2)}_{ll}$, $\rho^{(Z_2)}_{kl}$ and $\rho^{(Z_2)}_{lk}$ follow from (2.15) and (2.16) by the substitution

$$g_L D_{11} \longrightarrow g_R D_{22}\,, \qquad g_R D_{12} \longrightarrow g_L D_{21} \qquad \xi \longrightarrow -\xi\,, \tag{2.19}$$

respectively.

From expression (2.11) it follows that in this model the magnetic dipole μ_γ and the electric quadruple Q_γ moments of Z_1-boson have the same values as those of the SM Z-boson. However, the weak dipole μ_{Z_1} and the weak quadruple Q_{Z_1} moments of Z_1-boson (by this we mean the moments corresponding to the $W_1 W_1 Z_1$-couplings) are defined by the expressions

$$\mu_{Z_1} = (\mu_Z)_{\rm SM} + \frac{e_Z \Delta k_1}{2m_{W_1}}, \qquad Q_{Z_1} = (Q_Z)_{\rm SM} - \frac{e_Z \Delta k_1}{m^2_{W_1}}, \tag{2.20}$$

where $(\mu_Z)_{\rm SM}$ and $(Q_Z)_{\rm SM}$ are the multiple moments values in the SM, $\Delta k_1 = 2(\rho_{11}^{(Z_1)} - e_Z)$, $e_Z = e \operatorname{ctg} \theta_W$.

I should like to remind that in any LR model one cannot avoid the existence of neutrinos masses. The crucial question that one must answer in these models is why the neutrino masses are so much smaller than the corresponding charged lepton masses. There is, fortunately, a very nice explanation for this problem which ties the existence of low mass neutrinos with the possibility of large PV in the theory. The V–A approximate structure of the weak interactions, arising from the asymmetric breakdown in which $m_{W_R} \gg m_{W_L}$, has a counterpart in the neutrino mass spectrum. Namely, in the limit $m_{W_R} \longrightarrow \infty$ in which the weak interaction become purely left-handed the masses of the left-handed and right-handed neutrinos tend to 0 and ∞, respectively. That was shown in Ref. [2] with the help of the see-saw mechanism (SSM) proposed in Ref. [7]. It is easy to see that SSM can be used in the continuous LR model without any changes. So, following the results of Ref. [2], in (2.12) and (2.13) we must make the substitution

$$\left.\begin{array}{l} \nu_{l\rm L} \longrightarrow \nu_l \cos\alpha - N_l \sin\alpha \\ \nu_{l\rm R} \longrightarrow N_l \cos\alpha + \nu_l \sin\alpha \end{array}\right\}, \tag{2.21}$$

where ν_l and N_l are the Majorana spinors describing the mass eigenstates, and the mixing angle α is very small.

3. Physical implications

In this section we investigate some processes in the continuous LR model. First of all, we consider the decay width of Z_n-boson into fermion pairs. It is worth to remind that as a rule the cross sections of the process going in the second and higher orders of the perturbation theory are compared with the experiment. Such parameters of New Physics as the masses of the additional bosons m_{Z_2} and m_{W_2}, to which the cross sections are the most sensitive, enter the second order processes already at the tree level. Then varying m_{Z_2}, m_{W_2} and the values of the other parameters of both LR1 and LR2, one might find the set of parameters which could reproduce the SM in the neutral sector (see, for example, Refs [8, 9]). However, the first order processes going with the participation of the SM particles are less sensitive to the variation of m_{Z_2} and m_{W_2} because the effects connected with them enter the cross sections at the level of radiative corrections (RC) only. Thus the investigation of these processes allows one to reduce the class of the SM extensions corresponding to the experiment significantly. The decay of Z-boson is the example of such a process. In our investigation we shall be limited to taking into account the oblique RC *i.e.* RC caused by self-energy diagrams. As it is known they lead to the change of the scale of the Born approximation (BA), which is called improved AB after including RC. In this approximation the decay width of Z_n-boson into fermion pair is determined by the expression

$$\Gamma_{Z_n\to \mathrm{f\bar{f}}} = \frac{m_{Z_n}\rho N_C}{48\pi\sqrt{1+\beta^2}}\sqrt{\left(1-\frac{2m_\mathrm{f}}{m_{Z_n}}\right)^2} \times \left[\left(1+\frac{2m^2}{m^2_{Z_n}}\right)\left(g^\mathrm{f}_{Vn}\right)^2+\left(1-\frac{4m^2_\mathrm{f}}{m^2_{Z_n}}\right)\left(g^\mathrm{f}_{An}\right)^2\right], \quad (3.1)$$

where

$$\rho = 1+\Delta\rho_{Z_2}+\Delta\rho_{W_2}+\Delta\rho_t+\ldots, \qquad \Delta\rho_{Z_2}\simeq\left(\frac{m_{Z_2}\sin\phi}{m_{Z_1}}\right)^2,$$

$$\Delta\rho_{W_2}\simeq\left(\frac{m_{W_2}\sin\xi}{m_{W_1}}\right)^2, \qquad \Delta\rho_t\simeq\frac{3G_F m_t^2}{8\sqrt{2}\pi^2}, \qquad \beta=\frac{m_{W_1}}{m_{W_2}},$$

$$N_C=\begin{cases}1, & \text{for leptons,}\\ 3\left(1+\frac{\alpha_{QCD}}{\pi}\right), & \text{for quarks,}\end{cases}$$

$f\neq b$, m_f is the mass of the fermion f, the effective s^2_W is connected with the $s^2_W=\sin^2\theta_W$ of the SM by the relation

$$s^2_W=\bar{s}^2_W-\frac{\bar{s}^2_W\bar{c}^2_W}{(\bar{c}^2_W-\bar{s}^2_W)}(\Delta\rho_{Z_2}+\Delta\rho_{W_2}).$$

With dots in the definition ρ we marked contributions of Higgs particles, heavy right-handed Majorana neutrinos, *etc.* For f $= b$ in (3.1) the changes should be done [10]

$$\left.\begin{array}{l}1+\Delta\rho_t\to\frac{1-\Delta\rho_t}{3},\\ s^2_W\to s^2_W\frac{1+2\Delta\rho_t}{3}\end{array}\right\}. \quad (3.2)$$

Using (3.1) one could show the validity of the approximate relation

$$\sum_{f} \frac{\Gamma_{Z_1 \to f\bar{f}}}{\Gamma_{Z_2 \to f\bar{f}}} \simeq \frac{m_{Z_1}}{m_{Z_2}} . \tag{3.3}$$

Now we take into account the results of LEP I experiments to obtain the bound on ϕ for the CSLR, *i.e.* in the case $g_{\rm L} = g_{\rm R} = \frac{e}{\sin\theta_W}$ and $g' = \frac{e}{\sqrt{\cos 2\theta_W}}$. The experimental values of high energy quantities used in our analysis are: **(a)** the total decay width $\Gamma_Z = (2.487 \pm 0.01)$ GeV; **(b)** the partial width to charged leptons $\Gamma_e = (83.2 \pm 0.55)$ MeV; **(c)** the axial and the vector couplings $(g_V^e)^2 = (1.16 \pm 0.41) \times 10^{-3}$, $(g_A^e)^2 = 0.2493 \pm 0.0013$. Comparing the theoretical and experimental values of the quantities quoted above one can conclude that the allowed values of the angle φ are

$$\varphi = \varphi_0 \pm \Delta\varphi , \tag{3.4}$$

where $\varphi_0 = 0°$, $180°$ and $\Delta\varphi = 3° - 5°$. To obtain the precise answer a complete knowledge of all the RC is needed and we are going to do this elsewhere. It should be noted that the current constraints on parameters of LR models which include RC are found for LR2 only. However, the complete calculation RC even at one-loop level was not done elsewhere. Thus, in Ref. [9] the RC connected with the new gauge structure relatively the $SU(2)_L \times U(1)$ were ignored. The obtained constraints on parameters of LR2 which are found on the analysis of high-energy processes at LEP I and SLC connected with low-energy data (atomic parity violation, and νe, eD, μC scatterings) have the form

$$m_{W_2} > 477 \text{ GeV}, \ |\xi| < 0.031, \ 0.008 < \tan\phi < 0.044, \ m_{Z_2} > 564 \text{ GeV} . \tag{3.5}$$

In Ref. [11] at the analysis of Γ_Z it was not taken into account the RC due to the W_2-boson exchange and the angle ξ was considered to be 0. That led to the increase of the low bound on m_{Z_2} up to 950 (800) GeV at $m_H = 100$ GeV (1 TeV). It should be noted that in the case $\varphi = 180°$ we have the model which follows from LR2 by the change of the sign in the Z_n couplings to fermions and W_n-bosons only. Therefore, the analysis of Refs [9, 11] give the same bounds on the model parameters for the case $\varphi = 180°$.

For Z_2-bosons the decays into the channels

$$Z_2 \to W_i^- W_k^+ \tag{3.6}$$

should be kinetically allowed (i, $k = 1, 2$). The calculations of those widths lead to the result

$$\Gamma_{Z_2\to W_i^- W_k^+} = \frac{c_W^2\left(\rho_{ik}^{(Z_2)}\right)^2 m_{Z_2}\rho}{96\sqrt{(1+\beta^2)}}\sqrt{\left(1-\frac{M_{ik}}{y_{ik}}\right)^2-\left(\frac{2}{y_{ik}}\right)^2}$$
$$\times\left[2+\frac{y_{ik}^2}{4}-M_{ik}\left(\frac{19M_{ik}^2}{4}-2y_{ik}+\frac{32-5M_{ik}}{2y_{ik}}\right)\right] \qquad (3.7)$$

where

$$y_{ik}=\frac{m_{Z_2}^2}{m_{W_i}m_{W_k}}, \qquad M_{ik}=\frac{m_{W_i}^2+m_{W_k}^2}{m_{W_i}m_{W_k}}.$$

The decay width $\Gamma_{Z_2\to W_1^- W_1^+}$ displays a strong dependence on the angle φ. For example, in the case CSLR it increases by an order of magnitude when varying φ from 0° up to 5°. It should be noted that the Z_2-boson decay into $W_1^-W_2^+$ is suppressed relatively to the decay into $W_1^-W_1^+$ by the factor

$$\left(\frac{m_{W_1}\sin 2\xi}{m_{W_2}}\right)^2.$$

We now proceed to consider the single production of Z_n-boson at hadron colliders. In the lowest order of Drell–Yan approximation the process

$$ab \to Z_n + X\,, \qquad (3.8)$$

where a, b = p, $\bar{\text{p}}$, is caused by the subprocess

$$\text{q}_i\bar{\text{q}}_i \to Z_n \qquad (3.9)$$

(i is the quark flavour). Let us introduce the rapidity variable of Z_n-boson in the center of mass system of the reaction (3.8)

$$y=\tfrac{1}{2}\ln\left|\frac{E+\text{p}_{\|}}{E-\text{p}_{\|}}\right|,$$

where E and $\text{p}_{\|}$ are the energy and the Z_n-boson longitudinal momentum, respectively. Then the differential cross section can be written in the form

$$\frac{d\sigma}{dy}=\frac{\pi}{24m_{Z_n}^2\sqrt{(1+\beta^2)}}\sum_i\left[f_{\text{q}_i}^{(a)}(x_a,Q^2)f_{\bar{\text{q}}_i}^{(b)}(x_b,Q^2)+(\text{q}_i\leftrightarrow\bar{\text{q}}_i)\right]$$
$$\times\left[\left(g_{Vn}^{\text{q}_i}\right)^2+\left(g_{An}^{\text{q}_i}\right)^2\right], \qquad (3.10)$$

where

$$\tau = \frac{m^2_{Z_n}}{s}, \quad x_{a,b} = \sqrt{\tau}\exp(\pm y), \quad s = (\mathrm{p}_a + \mathrm{p}_b)^2, \quad f^{(a)}_{\mathrm{q}}(x_a, Q^2)$$

is the distribution function of the quark flavour i in hadron a, the parameter Q^2 whose value is of order $\hat{s} = \left(\mathrm{p}_{\mathrm{q}_i} + \mathrm{p}_{\bar{\mathrm{q}}_i}\right)^2$ includes QCD corrections in the leading logarithmic approximation.

Having integrated (3.10) we obtain the total cross section in the form

$$\sigma = \frac{1}{24 m^2_{Z_n}\sqrt{(1+\beta)^2}} \sum_i \left[\left(g^{\mathrm{q}_i}_{Vn}\right)^2 + \left(g^{\mathrm{q}_i}_{An}\right)^2\right] \tau \frac{dL_{\mathrm{q}_i\bar{\mathrm{q}}_i}}{d\tau}, \tag{3.11}$$

where the differential luminosity $\tau dL_{\mathrm{q}_i\bar{\mathrm{q}}_i}/d\tau$ is defined by the following expression

$$\tau \frac{dL_{\mathrm{q}_i\bar{\mathrm{q}}_i}}{d\tau} = \int\limits_{\tau}^{1} \left[f^{(a)}_{\mathrm{q}_i}(x, Q^2) f^{(b)}_{\bar{\mathrm{q}}_i}\left(\frac{\tau}{x}, Q^2\right) + (\mathrm{q}_i \leftrightarrow \bar{\mathrm{q}}_i)\right] \frac{dx}{x} \tag{3.12}$$

$(x = 2\sqrt{\tau} shy)$.

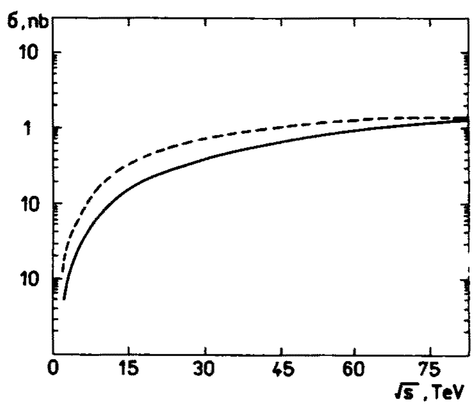

Fig. 1. The total cross section for Z_2-production in $\mathrm{p\bar{p}}$ (solid line) and pp (dashed line) collisions.

We shall neglect the distribution of the sea quarks c, s, t, b and use the parametrization of the parton distribution of Ref. [12] (set 2). In Fig. 1 it is shown the total cross section of the reaction (3.8) *versus* $\sqrt{s}$ with $n = 2$ for LR2 in case of pp (solid line) and $\mathrm{p\bar{p}}$ (dashed line) collisions. In our numerical calculations we used the following values of parameters

$$m_{Z_2} = 600 \text{ GeV}, \quad \sin^2\theta_W = 0.23, \quad \xi = 10^{-2}, \quad \phi = 4\times 10^{-2}. \tag{3.13}$$

The obtained results show that the advantage of $p\bar{p}$ collisions stretches to the field of larger values τ than it takes place at the Z-production in the SM. For the SM the advantage of the $p\bar{p}$ beams is important only when $\sqrt{\tau} \geq 0.01$, while for LR2 in (3.13) it takes place when $\sqrt{\tau} \geq 0.6$.

From (3.11) with $n = 1$ it follows that the values of the SM deviations for LR2 are of order $\sin^2\phi$. One could see that the $\sigma_{pp\to Z_2+X}, \sigma_{p\bar{p}\to Z_2+X}$ is smaller by a factor 10^{-2} than the $\sigma_{pp\to Z_1+X}, \sigma_{p\bar{p}\to Z_1+X}$. However, the number of Z_2-bosons produced in the reaction (3.8) is large enough. For example, the number of events in a year predicted by LR2 are the following: (a) for Tevatron ($\sqrt{s} = 2$ TeV, $\int Ldt = 10^3 pb^{-1}$) are 5×10^4; (b) for LHC ($\sqrt{s} = 17$ TeV, $\int Ldt = 10^4 pb^{-1}$) are 5×10^6; (c) for SSC ($\sqrt{s} = 40$ TeV, $\int Ldt = 10^4 pb^{-1}$) are 10^7. The leptonic decays

$$Z_2 \to e^-e^+, \mu^-\mu^+$$

should produce a clear signal with essentially no background expecting such instrumental problems as e/π separation only. The observation of reaction (3.8) also provides nice opportunity to study $\Gamma_{Z_2} \to W_1^- W_1^+$. It is known [13] that the Z_2-production and its subsequent decay into $W_1^- W_1^+$ lead to the "anomalous" events, *i.e.* to the events with large transverse momentum of the lepton pair recoiling against hadron jets. The relation of "anomalous" events to decays

$$Z_2 \to e^-e^+$$

greatly depend on φ and m_{Z_2}. Thus for example in CSLR at $m_{Z_2} = 600$ GeV

$$\Gamma_{Z_2\to e^-e^+} + \simeq 0.67 \text{ GeV} \qquad \Gamma_{Z_2\to W_1^- W_1^+} \simeq 1.99 \text{ GeV}$$

when $\varphi = 0$ and

$$\Gamma_{Z_2\to e^-e^+} \simeq 0.81 \text{ GeV} \qquad \Gamma_{Z_2\to W_1^- W_1^+} \simeq 19.57 \text{ GeV}$$

when $\varphi = 5°$. So the investigation of the momentum distribution of the lepton pairs in the final channel of the reaction (3.8) is a good tool for the angle φ definition. Because of a large value of the QCD background it is expedient to analyze the effective mass distribution of the final state of the reaction (3.8) at Z_2-peak (the ratio signal/background is maximum here).

Now we consider the reaction

$$f_j\bar{f}_j \to W_i^- W_k^+, \tag{3.14}$$

where $i, k = 1, 2$. Again, we shall be limited to including the RC at the level of improved BA. The differential and the total cross section of the reaction

(3.14) in this model follow from the corresponding formulae of the Ref. [14] after the multiplication on ρ and the changes

$$\alpha(0) \to \frac{\alpha(0)}{1-\Delta r}, \qquad \bar{s}_W^2 \to s_W^2, \qquad \frac{\rho_{kl}^{(n)}}{\sqrt{2}\sin\theta_W} \to \rho_{kl}^{(Z_n)},$$
$$\frac{-eV_n^f}{\sqrt{2}\sin\theta_W} \to g_{Vn}^f, \qquad \frac{-eP_n^f}{\sqrt{2}\sin\theta_W} \to g_{An}^f. \tag{3.15}$$

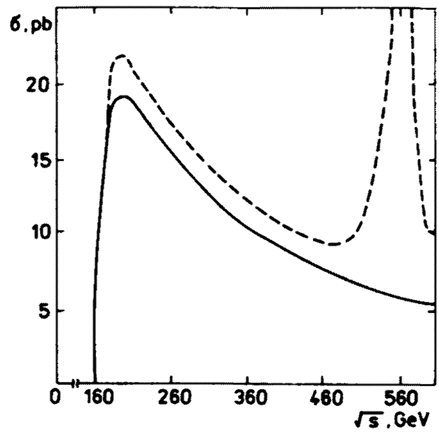

Fig. 2. The total cross section for the W-pair production in e^+e^- collisions. The solid and dashed lines stand for SM and LR2, respectively.

In Fig. 2 we present the total cross section of the reaction (3.14) with $i = k = 1$ and $f_j = e^-$ as a function of the energy in the center of mass system $\sqrt{2}$. The dashed line corresponds to the LR2, the solid one does to the SM. In numerical calculations we used the following parameters values

$$m_{Z_2} = 564 \text{ GeV}, \quad m_{W_2} = 477 \text{ GeV}, \quad m_{W_1} = 80.13 \text{ GeV},$$
$$m_{Z_1} = 91.172 \text{ GeV}, \; m_t = 145 \text{ GeV}, \; \Phi = 2\times 10^{-2}, \; \xi = 4 \times 10^{-2}. \tag{3.16}$$

The maximum value of the total cross section of the W-pair boson production in the LR2 has the order 21 *pb* at $\sqrt{2} \simeq 200$ GeV. That is larger than the total cross section of the fermion pair production which dominates below that energy.

In practice it is convenient for the analysis to use the quantity δ which is determined by the relation

$$\delta = \frac{(\sigma)_{\rm SM} - (\sigma)_{\rm LR}}{\sqrt{(\sigma)_{\rm SM}}}\sqrt{\rm LT}, \tag{3.17}$$

where LT is the integrated luminosity of the collider in units of pb^{-1}. The quantity δ gives the deviations from the SM expressed in the standard error units. In Fig. 3 the dependence of δ on $\sqrt{s}$ is shown for LR2 at the values of parameters (3.16) and LT= 100 pb^{-1}. As it follows from Fig. 3 the deviations from the SM reach the values of some standard errors already at the energies of the order of 200 GeV. Comparing the obtained results with those of Ref. [15] we see that the inclusion of RC leads to the sizable increase of the deviations from the SM for the LR2. It should be noted that the main contribution is caused by the redefinition of s_W, *i.e.* the RC connected with the self-energy diagrams of the neutral and the charged gauge bosons. The total cross section greatly depends on m_{Z_2} while the dependence on ξ and Φ is weak enough. For LR2 the value 564 GeV is the low bound for the Z_2-boson mass. At the increasing m_{Z_2} the Γ_{Z_2} grows quickly what leads to the increase of the $(\sigma)_{LR2}$ in the energy region up to Z_2-peak. This circumstance together with the growth of the contribution from RC cause the increase of δ. One can see that the investigation of reaction (3.14) on LEP II will be decisive for symmetric LR models at allowed values of their parameters.

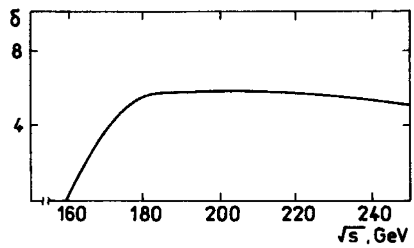

Fig. 3. The δ *versus* $\sqrt{s}$.

4. Conclusions

The continuous LR model is suggested to unify all the known LR models. This model is characterized by the orientation angle φ of $SU(2)_R$ generator in the group space. It should be stressed that the role of φ is not reduced to the redefinition of the neutral gauge bosons mixing angle ϕ. There is a principal difference between these quantities. The angle φ describes the rotation in three-dimensional space while the angle ϕ works in two-dimensional space only. The comparison of the symmetric version of this model with the experiments shows that the φ allowed values lie near 0° and 180°. It should be reminded that CSLR reproduces SM in the sector of SM particles at $\varphi = \xi = \phi = 0$. The case $\varphi = 180°$ follows from one $\varphi = 0°$

by the sign change in $L_{\rm NC}$ and L_{WWZ}, only. It has been shown that the cross sections of the reactions

$$Z_n \to {\rm f\bar{f}}\,, \quad {\rm p\bar{p}} \to Z_n + X\,, \quad {\rm f\bar{f}} \to {\rm W}_1^+{\rm W}_1^-$$

have the same values for $\varphi = 0°$ and $\varphi = 180°$.

The author would like to thank the organizers of the XXXII Cracow School of Theoretical Physics for constructive and nice atmosphere. He would also like to thank all the participants of the School for their responses and criticism.

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