PROBING PHYSICS BEYOND THE STANDARD MODEL*

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The two versions of the electroweak models are investigated: the left-right (LR) model (both the symmetric and asymmetric cases) and the composite one (CM). The total cross sections of the W-pair production are found in both models. The analysis of LR model results show that already at LEP II energies we can either establish or limit such parameters of LR model as $\Delta\rho_M$, m_{ν_R} and g_R. It is also found that the investigation of the single production of Z_2-boson in pp collisions will be the good tool for the g_R definition. For the CM the deviations value from the SM is mainly defined by the multiple moments values.

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1. Introduction

Experimental data, collected up to now, strongly support the electroweak model based on the $SU(2)_L \times U(1)_Y$ gauge group (SM) as being the correct description of physics at currently accessible energies. No doubt that the SM exhibits elements of truth. However, in spite of its impressive success, the SM leaves many fundamental questions unanswered: the motivation in the choice of the symmetry group, the Higgs sector shows, hierarchy problem, number of generation, origin of parity and CP violation, ... All that leads us to the belief that the SM is not a fundamental theory, but is only a low-energy approximation to a more general unified theory. The search for the solution of these problems has resulted in construction of the SM extensions which can be divided into two main categories: (a) the model having physical Higgs but evading the deficiencies of the SM, (b) the model without any physical Higgs. The first class

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consists of the superstring theories. If we take $E_8 \times E_8$ from phenomenological interest, it breaks by the loop mechanism down to some subgroup of $E_8 \times E_6$. The group E_6 is one of the most interesting candidates as unifying group by the following reasons: (a) it contains as subgroups the symmetry groups of the most popular grand unified theories SO(10), SU(5); (b) it is the only phenomenologically acceptable group which can arise from ten-dimensional superstring theories after compactification on the Calabi–Yau manifold. In the group E_6 the $G_{LR} = SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ and the $SU(2)_L \times U(1) \times U'(1)$ ($U''(1)$) electroweak symmetry groups can appear as the intermediate gauge structures. The composite models (CM's) belong to the second class. In such models the observed weak interactions are phenomenological manifestations of some basic confinement dynamics. The symmetry structure of the current-current Lagrangian of low energy weak interactions is obtained by imposing on the underlying Lagrangian either the $SU(2)_L$ [1] or the $SU(2)_L \times SU(2)_R$ [2] weak isospin symmetries. All these versions of the electroweak theories predict either new gauge bosons or the set of the excited states of the ordinary W, Z bosons. Therefore, if these new states, or the effects connected with them are found experimentally, we must have the criteria to single out the model corresponding to the experimental data.

This paper is devoted to the investigation of the models belonging to the both classes quoted above: (a) the symmetric and the asymmetric versions of the left-right (LR) model, (b) the composite models. The plan of the paper is the following: In Section 2 I assume the model which unifies the wide class of the different versions of the LR model. The obtained model is analyzed from the point of view of the experiments coming from LEP and future colliders. In Section 3 I consider the reaction

$$e^+e^- \longrightarrow W^+W^- \tag{1.1}$$

in CM's. I obtain the differential and the total cross sections in the analytical form and compare the results with those of the SM and the LR model. Finally, I summarize my work in Section 4.

2. LR models

A possible extension of the SM is the LR model based on the G_{LR} gauge group. This model accounts for many, but by no means all, physics problems which cannot get the satisfactory explanation within the SM. The parity violation (PV) in weak interactions is one of the examples of such problems. The observed near maximum PV in low energy weak interaction may be interpreted in LR models, as arising out of spontaneous breaking of parity and consequent nonvanishing neutrino masses (possibly required by astrophysics and cosmology). There are other important reasons for considering these models. For example, LR models give the comprehensive picture of the fermion spectrum [3] (from under 20 eV for the electron neutrino to over 100 GeV for the yet undiscovered top quark), the quantum numbers of group U(1) are identified with B–L (instead of Y having no physical meaning) what allows to link the breaking of parity and breaking of B–L, LR models allow for the generation of CP violation *via* the spontaneous symmetry breaking mechanism and can account for its strength by relating it to the suppression of right-handed currents [4].

There are a lot of papers in which such models are considered [5–8]. All of them could be unified into the one common model suggested in Refs. [9, 10]. In that model the neutral current interaction is parameterized by five parameters (instead of four ones as it takes place in ordinary LR model): two Z masses, the mixing angle of the neutral gauge bosons Φ, Weinberg angle θ_W, and the angle φ characterizing the orientation of the $SU(2)_R$ generator in the group space. Changing φ I can reproduce all the known LR models. In the case symmetric LR model ($g_L = g_R$) the analysis of Z_1 decay parameters being made at LEP I gives the following bounds on φ [9]

$$\varphi \leq few \times 10^{-2}\text{rad}\,.$$

One could show that the SM is reproduced in the SM particles sector if and only if the following conditions are fulfilled

$$\varphi = \Phi = \xi = 0, \quad g_L = g_R = e s_W^{-1}, \quad g' = e(c_W^2 - s_W^2)^{-1/2} \tag{2.1}$$

in the symmetric case and

$$\varphi = \Phi = \xi = 0,\ g_L = e s_W^{-1},\ g' = \pm(c_W^2 e^{-2} - g_R^{-2})^{-1/2},\ |g_R| > g_L s_W c_W^{-1} \tag{2.2}$$

in the asymmetric one (hereafter I use the notations of Ref. [9]). In conclusion, in this case, I pay some attention to the bounds of the g_R change in (2.2). If we start from the G_{LR} gauge group directly then the coupling constants g_L, g_R and g' are all arbitrary. Actually, they are not constants, but functions of $Q^2 = -p_\mu^2$ where p_μ is a typical momentum relevant to the process being considered. When the G_{LR} is embedded into a grand unified theory (GUT) then the three running coupling constants (CC's) g_L, g_R and g' all must come together at the grand unification scale which value should be consistent with the proton decay constraint. These demands impose the limitation on the choice of the unifying group. For example, the non-supersymmetric SU(N≤5) models are ruled out while the sypersymmetric SU(5) and SO(N≤10) models satisfy the demands above. The Q^2 dependence CC's can be calculated from the renormalization group equations in which these quantities enter quadratically. The choice both of the GUT and the scale of the underlying symmetry breaking fixes the absolute values of CC's at the electroweak scale (see, for example, the general LR model [11] in which $\varphi = 0$, $g_L = e s_W^{-1}$, $g' = (c_W^2 e^{-2} - g_R^{-2})^{-1/2}$, $g_L^2/2 \geq g_R^2 \geq g_L^2$). Up to now we cannot give preference to the definite GUT with confidence. Therefore, investigating the G_{LR} gauge theory as a low energy approximation of GUT we should consider all the possible g_R variation in (2.2). It should be remembered that the sign with which CC's enter the pieces of the Lagrangian are fixed by the choice of the signs in front of the fermion fields placed in the multiplets representations of GUT.

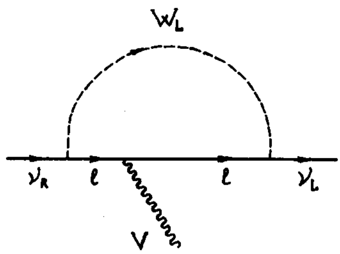

Fig. 1. The diagram corresponding to the decay $\nu_R \longrightarrow \nu_L + V$ ($V = \gamma,\ Z_{1,2}$).

Since I am going to take into account the effects connected with the existence of the heavy right-handed neutrino ν_R then I briefly discuss some of its properties. The value of the mass of the ν_R depends on two parameters: (a) the Yukawa coupling of the right-handed Higgs triplets Δ_R and (b) the mass of the right-handed W_R boson. The bounds on the m_{ν_R} could be obtained from the cosmological considerations as well as weak decay processes. For example, if ν_R are Majorana particles then the known bound on the neutrinoless double β-decay half-life of ^{76}Ge gives the lower bound [12]

$$m_{\nu_R} > 63 \text{ GeV} \left(\frac{1.6 \text{ TeV}}{m_{W_R}}\right)^4 .$$

The upper bound obtained in Ref. [11] on m_{ν_R} practically does not depends on m_{W_R} and is about 1000 GeV.

If the right-handed neutrino is too heavy so that the condition

$$m_{\nu_R} > m_{W_i} + m_l \,, \tag{2.3}$$

where m_l is the mass of a charged lepton, takes place ($i = 1, 2$) then it would decay into the channels

$$\nu_R \longrightarrow W_i^{\pm} + l^{\mp} . \tag{2.4}$$

Their decay widths are

$$\Gamma_{\nu_R} \longrightarrow W_{i+l} = \frac{g_R^2 h_i m_{\nu_R}}{32\pi}\left(y_i + \frac{2}{y_i^2} - \frac{3}{y_i}\right), \tag{2.5}$$

where $y_i = (m_{\nu_R}/m_{W_i})^2 \qquad h_i = \begin{cases} \sin^2\xi, & i=1 \\ \cos^2\xi, & i=2 \end{cases}.$

If the neutrino is the Dirac particle then we have the additional factor 1/2 in the right-hand side of Eq. (2.5).

The decays

$$\nu_R \longrightarrow \nu_L + V \qquad (V = \gamma,\ Z_{1,2}) \tag{2.6}$$

also could be allowed. These processes proceed *via* the Feynman graph drawn in Fig. 1. The decays widths are defined by the expression

$$\Gamma_{\nu_R \to \nu_L + V} = \frac{m_{\nu_R}^5}{2\pi}\left(\sum_{i=1} f_i\right)\tau_v, \tag{2.7}$$

where

$$f_i = \frac{g_L g_R}{96\pi^2 m_{W_i}^2}\sin\xi\left(\frac{m_l}{m_{W_i}}\right)^2,$$

$$\tau_v = \begin{cases} \frac{e^2}{2} & V = \gamma \\ \frac{1}{6}[(g_{Vn}^l)^2 + (g_{An}^l)^2]\sqrt{1-(4m_{Z_n}/m_{\nu_R})^2} & V = Z_n \end{cases}.$$

However, the $\Gamma_{\nu_R \to \nu_L + V}$ are very small. For instance, at $m_{\nu_R} = 500$ GeV, $m_l = 1$ GeV and $g_L = g_R = e/s_W$ we have

$$\Gamma_{\nu_R \to \nu_L + \gamma} \simeq 7.52\ \sin^2\xi\ (\text{GeV}).$$

Analogously, the channel $\nu_R \longrightarrow W_1 + 1$ is very suppressed by the factor $\sin^2\xi$. Therefore, for the case $m_{W_R} > m_{\nu_R}$ which I shall, in what follows, consider I neglect $\Gamma_{\nu_R \to \text{all}}$ for the sake of simplicity.

Now I consider the reaction

$$f_i \bar{f}_i \longrightarrow W_k^- W_n^+, \tag{2.8}$$

where i is the flavour of a fermion and $k,\ n = 1, 2$. I shall be limited by including the RC at the level of improved Born approximation [13]. I assume that neutrinos are Dirac particles. Then the total cross section for the polarized initial and unpolarized final particles is defined by

$$\sigma^{(kn)} =$$

$$\frac{\beta\rho}{56\pi s}\Big\{\Big[\frac{1-\lambda\bar{\lambda}}{4}\Big(|Q_i e^2[1+(-1)^{k+n}]+\sum_{l=1}^{2}2\rho_{kn}^{(z_l)}g_{Vl}^i s d_{z_l}|^2$$

$$+\Big|\sum_{l=1}^{2}2\rho_{kn}^{(z_l)}g_{Al}^i s d_{z_l}\Big|^2\Big]\Big)+\frac{\lambda-\bar{\lambda}}{4}\Big(4Q_i e^2[1+(-1)^{k+n}]$$

$$\times\sum_{l=1}^{2}\rho_{kn}^{(z_l)}g_{Al}^i Re(d_{z_l})s+4\sum_{l=1}^{2}\rho_{kn}^{(z_l)}g_{Al}^i d_{z_l}s\sum_{j=1}^{2}\rho_{kn}^{(z_j)}g_{Vj}^i d_{z_j}^* s\Big)\Big]$$

$$\times D_1^{(kn)}+\frac{(1-\bar{\lambda})(1+\lambda)}{4}g_{\rm L}^2 a_+^{kn}\Big(Q_i e^2+\sum_{l=1}^{2}\rho_{kn}^{(z_l)}(g_{Vl}^i+g_{Al}^i)sRe(d_{z_l})\Big)$$

$$\times D_2^{(kn)}+\frac{(1+\bar{\lambda})(1-\lambda)}{4}g_{\rm R}^2 a_-^{kn}\Big(Q_i e^2+\sum_{l=1}^{2}\rho_{kn}^{(z_l)}(g_{Vl}^i-g_{Al}^i)sRe(d_{z_l})\Big)$$

$$\times D_{2\nu}^{(kn)}+\frac{(1+\bar{\lambda})(1-\lambda)}{4}g_{\rm R}^4 b_+^{kn}D_{3\nu}^{(kn)}+\frac{(1-\bar{\lambda})(1+\lambda)}{4}g_{\rm L}^4 b_-^{kn}D_3^{(kn)}\Big\}, \tag{2.9}$$

where $s=(p_{W_n}+p_{W_k})^2$, $m_{W_{k,n}}\equiv m_{k,n}$.

$$\beta=\Big[\Big(1-\frac{m_n^2+m_k^2}{s}\Big)^2-\Big(\frac{2m_n m_k}{s}\Big)^2\Big]^{1/2},\ d_{z_l}=(s-m_{z_l}^2+im\Gamma_{z_l})^{-1},$$

$$a_\pm^{kn}=\frac{1+(-1)^{k+n}\mp[(-1)^k+(-1)^n]\cos 2\xi\pm[1-(-1)^{k+n}]\sin 2\xi}{4},$$

$$b_\pm^{kn}=\frac{1+(-1)^{k+n}-(-1)^{k+n}\sin^2 2\xi\pm[(-1)^k+(-1)^n]\cos 2\xi}{4},$$

$$D_1^{(kn)}=\beta^2\frac{(s-m_n^2-m_k^2)^2+12s(m_n^2+m_k^2)+8m_n^2m_k^2}{12m_n^2m_k^2},$$

$$D_{2v}^{(kn)}=\frac{(s-m_n^2-m_k^2)}{m_n^2m_k^2}\Big\{2(m_n^2+m_k^2)+\frac{s\beta^2}{6}+\frac{2m_n^2m_k^2}{s}-\frac{m_{\nu_{\rm R}}^4}{2s}$$

$$\times\Big[2m_{\nu_{\rm R}}^2-m_n^2-m_k^2+s+\frac{8m_n^2m_k^2}{m_n^2+m_k^2-s}\Big]\Big\}$$

$$+\frac{4L}{\beta s}\left[-m_n^2-m_k^2-\frac{m_n^2m_k^2}{s}+m_{\nu_R}^2(m_n^2+m_k^2-s)\right.$$
$$\left.\times\left(\frac{m_n^2+m_k^2}{2m_n^2m_k^2}+\frac{3}{4s}-\frac{m_{\nu_R}^4}{4sm_n^2m_k^2}\right)+\frac{\beta^2 sm_{\nu_R}^4}{4m_n^2m_k^2}\right],$$

$$D_{3\nu}^{(kn)}=\frac{1}{12m_n^2m_k^2}[s^2\beta^2-24m_n^2m_k^2+12s(m_n^2+m_k^2)$$
$$-6m_{\nu_R}^2(s-m_k^2-m_n^2+3m_{\nu_R}^2)]+C_\nu\left\{[m_{\nu_R}^2(m_n^2+m_k^2-s-m_{\nu_R}^2)\right.$$
$$\left.-m_n^2m_k^2]\left(2+\frac{m_{\nu_R}^4}{2m_n^2m_k^2}\right)+\frac{sm_{\nu_R}^4(m_n^2+m_k^2)}{m_n^2m_k^2}\right\}$$
$$+\frac{L}{\beta s}\left\{2(s-m_n^2-m_k^2)+5m_{\nu_R}^2+\frac{m_{\nu_R}^2}{2m_n^2m_k^2}\right.$$
$$\left.\times[m_{\nu_R}^2(4m_{\nu_R}^2-3m_n^2-3m_k^2+3s)-4s(m_n^2+m_k^2)]\right\},$$

$$D_2^{(kn)}=D_{2\nu}^{(kn)}|_{m_{\nu_R}=0},\qquad D_3^{(kn)}=D_{3\nu}^{(kn)}|_{m_{\nu_R}=0},$$
$$C_\nu=[m_n^2m_k^2-m_{\nu_R}^2(m_n^2+m_k^2-m_{\nu_R}^2-s)]^{-1},$$

$$L=\ln\left|\frac{m_n^2+m_k^2+2m_{\nu_R}^2+\beta s-s}{m_n^2+m_k^2+2m_{\nu_R}^2-\beta s-s}\right|\qquad Q_i=\begin{cases}-1 & \text{charged leptons}\\ 2/3 & \text{u - like quarks}\\ 1/3 & \text{d - like quarks}\end{cases}.$$

The quantity ρ entering (2.9) appears due to the RC cased by all heavy particles and it is defined by

$$\rho=1+\Delta\rho_t+\Delta\rho_{\rm M}+\dots,\tag{2.10}$$

where $\Delta\rho_t\simeq 3G_Fm_t^2/8\sqrt{2\pi}$, m_t is the top quark mass, $\Delta\rho_{\rm M}$ arises due to the mixing in the gauge vector bosons sector W and has the general asymptotic form [13]

$$\Delta\rho_{\rm M}=c_0^2\left(\frac{m_{z_1}}{m_{z_2}}\right)^2-c_1^2\left(\frac{m_{W_1}}{m_{W_2}}\right)^2,\tag{2.11}$$

with c_0 and c_1 constant depending on the VEV's Higgs particles and the CC's. We should remember that when $\Delta\rho_{\rm M}$ is not equal to zero then the effective s_W^2 in the LR model is connected with the $\bar{s}_W^2=\sin^2\theta_W$ of the SM by the relation

$$s_W^2=\bar{s}_W^2-\frac{\bar{s}_W^2\bar{c}_W^2}{(\bar{c}_W^2-\bar{s}_W^2)}\Delta\rho_{\rm M}.$$

Additional contributions to ρ (dots in the ρ definition) come from Higgs particles, heavy right-handed neutrinos, *etc.* They could be both positive and negative. For example, the contribution from the standard Higgs boson is logarithmic in the Higgs mass and for $m_H \gg m_{W_1}$ it has the negative sign. The contribution connected with the presence of weak isotriplets Higgs bosons is [14]

$$(\Delta\rho_{\mathrm{TH}})_{\mathrm{L,R}} = -\frac{2\Delta_{\mathrm{L,R}}^2}{\nu^2},$$

where $\nu = 246$ GeV is the standard Higgs-doublets vacuum expectation value. I shall, in what follows, make the simplifying assumption that all these additional contributions to ρ are mutually cancelled. The constraint on ρ coming from the measurements of m_z at LEP and m_W/m_z at UA2 and CDF has the form

$$\rho = 1.0066 \pm 0.0058\,. \tag{2.12}$$

So, the upper and lower bounds on $\Delta\rho_{\mathrm{M}}$ coming from the experiments are

$$(\Delta\rho_{\mathrm{M}})_{\mathrm{upper,lower}} = (\Delta\rho)_{\mathrm{upper,lower}} - \Delta\rho_t - 1\,. \tag{2.13}$$

From the expression for $\sigma^{(kn)}$ we can see that its partial contributions $\sigma_{ij}^{(kn)}$ ($i,j = \gamma$, $\mathrm{Z}_{1,2}$, $\nu_{\mathrm{L,R}}$) fastly increase with the energy. Consequently, each $\sigma_{ij}^{(kn)}$ by itself would violate unitarity. However, the $\sigma^{(kn)}$ resulting from the sum of those contributions decreases with the energy according to

$$\sigma^{(kn)} \propto \frac{\ln s}{s} \tag{2.14}$$

due to a delicate gauge cancellation among $\sigma_{ij}^{(kn)}$ so that the unitarity bound is not a problem.

I start the analysis of the total cross section with the case of unpolarized electron–positron beams. For the comparison with the SM it is convenient to introduce the quantity δ which characterizes the experimental sensibility to the deviation from the SM

$$\delta = \frac{(\sigma)_{\mathrm{LR}} - (\sigma)_{\mathrm{SM}}}{\sqrt{(\sigma)_{\mathrm{SM}}}}\sqrt{\mathrm{LT}}\,, \tag{2.15}$$

where LT is the integrated luminosity of the collider in units of pb^{-1}, $(\sigma)_{\mathrm{LR}}$ and $(\sigma)_{\mathrm{SM}}$ are the total cross section summed on initial particles polarizations in the LR and standard models, respectively. The δ is an observable of the effect from new physics and it gives the deviations from the SM expressed in the standard error units. Previous work on this problem [15] has shown that in the LEP II energies region the total cross section of the LR model calculated at the tree level differs from that of the SM on the values of order of a few$\times 10^{-3}$. It is well known that the main contribution caused by the RC is connected with the redefinition of $\bar{s}_W$, *i.e.* with the quantity $\Delta\rho_M$. It should be reminded that because of the structure of $\Delta\rho_M$ given by Eq. (2.11), $\Delta\rho_M$ can have not only the positive and the negative values but also the zero ones. In order to get an idea about the values of deviations caused by RC I shall use the upper, lower and zero bounds of $\Delta\rho_M$ in my analysis of the obtained cross sections.

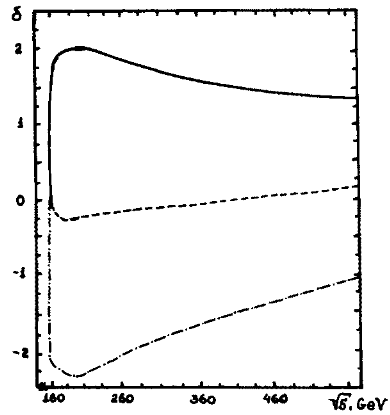

Fig. 2. The δ for the symmetric LR model as a function of $\sqrt{s}$ for $(\Delta\rho_M)_{\text{upper}}$ (solid line), $(\Delta\rho_M)_{\text{lower}}$ (dash-dotted line), and $\Delta\rho_M = 0$ (dashed line).

First I consider the case of the symmetric LR model. In Fig. 2 it is shown the δ of the reaction (2.8) with $i = k = 1$ and $f_j = e^-$ as a function of the energy in the center of mass system $\sqrt{s}$ for $(\Delta\rho_M)_{\text{upper}}$, $(\Delta\rho_M)_{\text{lower}}$, and $(\Delta\rho_M) = 0$ at LT=500 pb^{-1}. I note that at the given values of m_t $(\Delta\rho_M)_{\text{lower}}$ is negative. In numerical calculations I used the following values of the SP's [16, 17]

$$\varphi = 0 \quad m_{Z_2} = 800 \text{ GeV}, \quad m_{W_2} = 477 \text{ GeV}, \quad m_{\nu_R} = 400 \text{ Gev},$$

$$\Phi = 9.6 \times 10^{-3}, \quad \xi = 3.1 \times 10^{-2}. \tag{2.16}$$

From Fig. 2 we see that the possible deviations of the LR model from the SM lie within the region restricted by the curves $(\Delta\rho_M)_{\text{upper}}$ and $(\Delta\rho_M)_{\text{lower}}$. For the reaction

$$e^-e^+ \longrightarrow W_1^-W_1^+$$

δ displays a weak dependence on m_{ν_R} because of the $W_1\nu_R e$ coupling is proportional to $\sin\xi$. If we vary m_{ν_R} from 400 to 100 GeV then $|\delta|$ increases in value an order of 1%. The role of ν_R becomes essential for the processes

$$e^-e^+ \longrightarrow W_1^-W_2^+,\ W_2^-W_2^+$$

by virtue the fact that the $W_2\nu_R e$ coupling is proportional to $\cos\xi$.

We see that at the chosen value Φ and $\Delta\rho_M > 0$, $\delta(\sqrt{s})$ has the minimum in the energy region less than m_{Z_2}. The presence or the absence of this minimum is defined by the Φ value only. At the big values of Φ (4×10^{-2}) the minimum is absent and $\delta(\sqrt{s})$ monotonously increases to $\sqrt{s} = m_{Z_2}$ [15]. The analysis shows that the minimum position $(\sqrt{s})_{min}$ and its value δ_{min} significantly depend on the chosen values m_{Z_2} and Φ, in contrast, the change of the other SP's weakly influences both on δ_{min} and $(\sqrt{s})_{min}$. At $\Delta\rho_M \leq 0$ the function $\delta(\sqrt{s})$ crosses the axis of abscissae in the point whose position mainly depends on the mixing angles and m_{Z_2}. Referring to Fig. 2, the important feature of $\delta(\sqrt{s})$ in the LEP II energies region is the fast increase of its module with the growth of $|\Delta\rho_M|$. Therefore, in depending on the measurement precision of the total cross section one could either establish or limit the $\Delta\rho_M$ value.

The analysis shows that in the case of the asymmetric LR model the unpolarized total cross section displays a weak dependence on g_R.

Now I consider the case of the polarized electron–positron beams. Using the total cross section I can define the polarization asymmetry

$$A_{LR} = \frac{(\sigma)_{1,-1} - (\sigma)_{-1,1}}{(\sigma)_{1,-1} + (\sigma)_{-1,1}}, \tag{2.17}$$

where $(\sigma)_{\lambda,\bar{\lambda}}$ is the total cross section either of LR or standard models in the case when the electron and the positron have the helisities λ and $\bar{\lambda}$, respectively. In Fig. 3 I display A_{LR} as a function of $\sqrt{s}$ for the case of symmetric LR model. The dash-dotted line stands for $(\Delta\rho_M)_{upper}$, the dotted one does for $(\Delta\rho_M)_{lower}$ and solid one stands for $(\Delta\rho_M) = 0$ (the values of SP are defined by (2.16)). I also display, for comparison, the results of the SM (dashed line). Again we see that the deviation from the SM depend on $\Delta\rho_M$. The analysis also shows the weak sensibility A_{LR} to the variations of g_R and m_{ν_R}.

For the case completely right-polarized electrons $(\lambda = 1)$ and left-polarized positrons $(\bar{\lambda} = -1)$ there is a crucial difference in the total cross section behaviour for the LR and the standard models. The SM cross section is defined by the diagrams involving the s-channel exchanged bosons γ and Z, only. In the LR model there are the t-channel contributions connected with the exchange of the right-handed neutrino. Though, in this case the

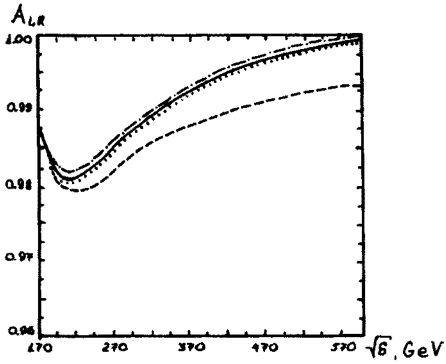

Fig. 3. The A_{LR} as a function of $\sqrt{s}$ for $(\Delta\rho_M)_{upper}$ (dash-dotted line), for $(\Delta\rho_M)_{lower}$ (dotted line) and for $(\Delta\rho_M) = 0$ (solid line). The results of the SM are represented by the dashed line.

total cross section at LEP II energies is two order of magnitude lower than the unpolarized total cross section its investigation would play the decisive role for the existence of the LR models. Let us define the function δ_{lr} by

$$\delta_{lr}(s) = \frac{(\sigma_{LR})_{1,-1} - (\sigma_{SM})_{1,-1}}{\sqrt{(\sigma_{SM})_{1,-1}}}\sqrt{LT}\,. \tag{2.18}$$

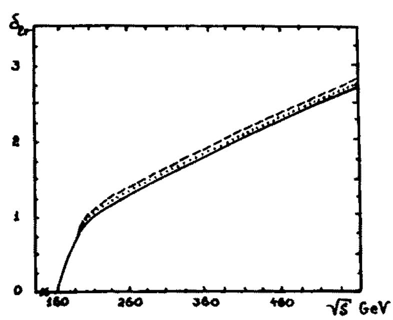

Fig. 4. The δ_{lr} *versus* $\sqrt{s}$ for the cases: (a) $m_{\nu_R} = 100$ GeV and $(\Delta\rho_M)_{upper}$ (solid line), (b) $m_{\nu_R} = 400$ GeV and $(\Delta\rho_M)_{upper}$ (dashed line), (c) $m_{\nu_R} = 100$ GeV and $(\Delta\rho_M)_{lower}$ (dotted line).

In Fig. 4 δ_{lr} as a function of $\sqrt{s}$ is presented for the following cases: (a) $m_{\nu_R} = 100$ GeV and $(\Delta\rho_M)_{upper}$ (solid line), (b) $m_{\nu_R} = 100$ GeV and $(\Delta\rho_M)_{lower}$ (dotted line), (c) $m_{\nu_R} = 400$ GeV and $(\Delta\rho_M)_{upper}$ (dashed line) (hereafter the other SP of the LR model are the same as in the case of Fig. 2). It should be noted here that δ_{lr} is also sensitive to the values of the mixing angles ξ and Φ. However, the dependence of δ_{lr} is just the same as δ and $A_{\rm LR}$ on Φ is much more stronger than one on ξ. For example, at $\Delta\rho_{\rm M} = 0$ the variation of Φ from 9.6×10^{-3} up to 9.6×10^{-4} yields the decrease of δ_{l_r} on one order of magnitude while the one of ξ from 3.1×10^{-2} up to 3.1×10^{-3} does the increase of δ_{l_r} on a few percents. It is immediately apparent that δ_{l_r} could be used for the $g_{\rm R}$ definition. In Fig. 5 I display δ_{l_r} *versus* $g_{\rm R}/g_{\rm L}$ at $\sqrt{s} = 196$ GeV.

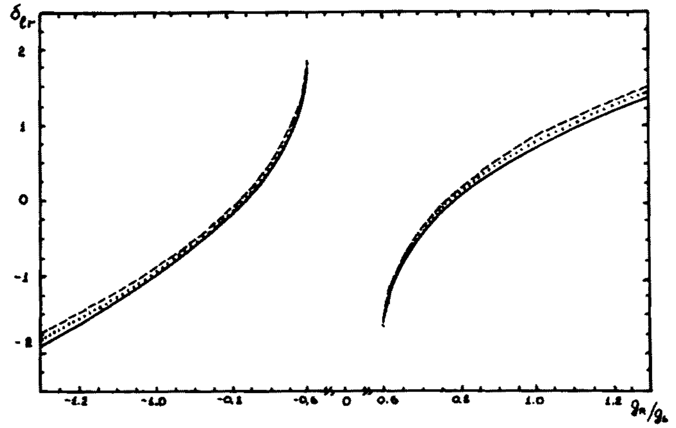

Fig. 5. The δ_{l_r} as a function of $g_{\rm R}/g_{\rm L}$ at $\sqrt{s} = 196$ GeV. The dashed (dotted) line stands for $(\Delta\rho_{\rm M})_{\rm upper}$ $((\Delta\rho_{\rm M})_{\rm lower})$ and $m_{\nu_{\rm R}} = 100$ GeV. The solid line stands for $(\Delta\rho_{\rm M})_{\rm upper}$ and $m_{\nu_{\rm R}} = 400$ GeV.

The hadron colliders provide the nice opportunity for the investigation of the extended gauge models. For example, the extra neutral gauge boson discovery limits at the LHC range from 2-3.5 TeV for an integrated luminosity of 10^4 pb^{-1} up to 4-5.5 TeV for an integrated luminosity of 5×10^5 pb^{-1}. Let us consider the possibilities of the hadron colliders for the $g_{\rm R}$ definition. It appears that the single production of $\rm Z_2$-boson is a good tool for this purpose. It adopt the spirit of the parton model. Then in the lowest order of Drell–Yan approximation the total cross section of the process

$$ab \longrightarrow {\rm Z}_n + \text{anything}, \tag{2.19}$$

where $a, b = p, \overset{(-)}{p}$, is defined by the expression

$$\sigma = \frac{1}{24 m^2_{\rm Z_2}\sqrt{1+\beta^2}} \sum_i \left[(g^{q_i}_{V_2})^2 + (g^{q_i}_{A2})^2\right] \tau \frac{dL_{q_i q_i}}{d\tau}, \tag{2.20}$$

where $\tau = \frac{m^2_{\rm Z_n}}{s}$, $s = (p_a + p_b)^2$, $f^{(a)}_{q_i}(x, Q^2)$ is the distribution function of the quark flavour i in hadron a, the parameter Q^2 which value is of order $\hat{s} = (p_{q_i} + p_{\bar{q}_i})^2$ includes QCD corrections in the leading logarithmic approximation and the differential luminosity $\tau d\mathrm{L}_{q_i\bar{q}_i}/d\tau$ is defined by the following expression

$$\tau \frac{d\mathrm{L}_{q_i\bar{q}_i}}{d\tau} = \int_{\tau}^{1} \left[f_{q_i}^{(a)}(x, Q^2) f_{\bar{q}_i}^{(b)}\left(\frac{\tau}{x}, Q^2\right) + (q_i \leftrightarrow \bar{q}_i) \right] \frac{dx}{x} . \qquad (2.21)$$

Again I shall use the quantity δ which is now determined by the relation

$$\delta = \frac{(\sigma \mathrm{Br})_{\mathrm{ALR}} - (\sigma \mathrm{Br})_{\mathrm{SLR}}}{\sqrt{(\sigma \mathrm{Br})_{\mathrm{SLR}}}} \sqrt{\mathrm{LT}} , \qquad (2.22)$$

where Br is the Z_2 branching ratios. At present δ gives the deviations from the symmetric LR model expressed in the standard error units. In order to eliminate backgrounds I shall only consider the following leptonic modes of Z_2-decay

$$Z_n \longrightarrow e^- e^+ , \ \mu^- \mu^+ . \qquad (2.23)$$

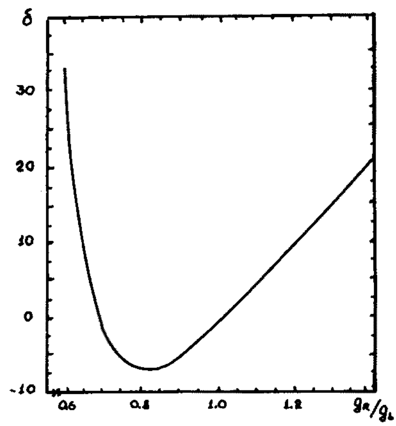

Fig. 6. The δ *versus* g_R/g_L for the pp collisions at LHC.

In Fig. 6 it is shown δ of the reaction (2.23) *versus* g_R/g_L for LHC ($\sqrt{s}$ = 17 TeV, LT= $10^4\, pb^{-1}$) in case of pp collidings. In my numerical calculation I neglected the distribution of the sea quarks c,s,t,b and used the parametrization of the parton distribution of Ref. [18] (section 2). I also used the following values of parameters

$$m_{Z_2} = 600 \text{ GeV} , \quad \xi = 3.1 \times 10^{-2} , \quad \phi = 9.6 \times 10^{-3} . \qquad (2.24)$$

3. The composite models

Now I shall consider the composite models. The basic assumption that underlies almost all composite model building is that the fundamental constituents often called preons interact by means of a confining hypercolor gauge interaction. According to current theoretical ideas, the non-abelian hypercolour theory should be asymptotically free, infrared confining and, of course, renormalizable. Below a characteristic energy scale Λ, the hypercolour interaction becomes strong and binds the preons into hypercolour–singlet states much in the same way as the conventional strong interactions among composite hadrons at $\Lambda_{\rm QCD}$. When the subenergies $\sqrt{s}$ exceed Λ, the manifestations of compositeness are very direct. At these energies, multiple production processes would dominate over the familiar two-body scattering processes. Examples of the sort of inelastic processes that may occur for a $f_i\bar{f}_i$ initial state are

$$f_i\bar{f}_i \longrightarrow \bar{f}_i f_i \bar{f}_i f_i, \;\; \bar{f}_i f_i \bar{f}_j f_j \;\; \bar{f}_j f_j \bar{f}_j f_j, \;\; W^-W^+W^-W^+ , \; etc. \tag{3.1}$$

The consequences of the compositeness are more subtle when $\sqrt{s}$ is small compared to Λ. Therefore, I focus on the signals which will be prominent for $\sqrt{s} < \Lambda$.

The classical test for substructure is to search for form factor effects, that is, deviations from the expected point-like behaviour in propagators and vertices. Many other tests of compositeness relies on the existence in CM's of new effective four-fermion contact interactions. Also the popular test is based on the investigation of the triple-vector boson couplings (TBC's). It is well known, that in CM's the TBC's are not fixed by the choice of the symmetry group as it happens in the SM. In the following I shall assume that the theory exhibits an $SU(2)_L\times U(1)$ global gauge symmetry and the exotic-fermion contact interactions are absent. So, the charged and the neutral currents have just the same forms as those in the SM. I shall also assume the P, C and T-invariance of the theory. Then the Lagrangian describing TBC's has the form [19]

$$\begin{aligned}\mathcal{L}_{WWV} =& ie_V\{(W^*_\sigma W_{\mu\sigma} - W_\sigma W^*_{\mu\sigma} + \partial_\sigma[k_v(W^*_\mu W_\sigma - W^*_\sigma W_\mu) \\ & - \lambda_V m_W^{-2}(W^*_{\mu\tau}W_{\sigma\tau} - W^*_{\sigma\tau}W_{\mu\tau})]\}V_\mu . \end{aligned} \tag{3.2}$$

The k_V and λ_V parameters define the magnetic [weak] dipole moment

$$\mu_\gamma = e\frac{1+k_\gamma+\lambda_\gamma}{2m_W}\left[\mu_Z = e_Z\left(\frac{1+k_Z+\lambda_Z}{2m_W}\right)\right], \tag{3.3}$$

and electric [weak] quadruple moment

$$Q_\gamma = e\frac{\lambda_\gamma - k_\gamma}{m_W^2} \qquad \left[Q_Z = e_Z\left(\frac{\lambda_Z - k_Z}{m_W^2}\right)\right]. \tag{3.4}$$

In the SM and the EGM's we have

$$k_\gamma = k_Z = 1\,, \qquad \lambda_\gamma = \lambda_Z = 0 \tag{3.5}$$

at the tree level. The RC increase the value of both the $k_{Z,\gamma}$ and the $\lambda_{Z,\gamma}$. The upper boundary of $\Delta k_\gamma = k_\gamma - 1 = 1.5\%$ and $\Delta\lambda_\gamma = 0.25\%$ has been obtained for the SM [20], for favourable values of the t-quark and the Higgs-boson masses. The introduction of either the extra fermion generation or the extra Higgs doublet (HD) increases the values of the $k_{Z,\gamma}$ and the $\lambda_{Z,\gamma}$ at the tree level. For instance, in the case of the two HD we have [20]

$$\Delta k_\gamma = 0.1\%\,, \qquad \Delta\lambda_\gamma = 0.03\%\,.$$

Each extra fermion SU(2) doublet gives a contribution to the Δk_γ less than 0.4%. Thus we see that the substantial deviation of the multiple moments (MM's) from their values in the SM can be obtained by a ridiculously large number of Higgs bosons or fermion families.

The electromagnetic MM's can be tested in the reaction

$$\mathrm{f}_i\bar{\mathrm{f}}_j \longrightarrow \mathrm{W}^-\gamma\,, \tag{3.6}$$

where f_i stands for one of the fermions $\nu, \mathrm{e}, \mathrm{u}, \mathrm{d}$, *etc.*, and $\bar{\mathrm{f}}_j$ for the correlated antifermions. It is known [21, 22] that the angular distribution for (3.6) vanishes at a certain angle provided the electromagnetic MM's have the values imposed by the SM. This critical angle, in the given frame, depends only on the charges of the particles and not on the masses or energy. In the case with the quarks when they are embedded into hadrons these radiation amplitude zeros (RAZ) manifest themselves as strong dips in angular distributions for the processes [22]

$$p\ p \longrightarrow \mathrm{W}^- + \text{ anything}\,. \tag{3.7}$$

The RAZ are also in the crossed channels of the reaction (3.6). All this is also true for the extended gauge models (EGM's). In the non-gauge theories with arbitrary values of the MM's, *e.g.*, the CM's the RAZ are absent. A good test for the weak MM's will be provided by the process

$$\mathrm{f}_i\bar{\mathrm{f}}_i \longrightarrow \mathrm{W}^-\mathrm{W}^+\,. \tag{3.8}$$

If I assume that the particles in the initial and final states are unpolarized then for the total cross section of this process I obtain the expression

$$\sigma = \frac{\alpha^2\beta\pi}{16x^2s}\Big(\sum_{i,j} M_{ij}\Big)\,, \tag{3.9}$$

where

$$i,j = \mathrm{Z},\gamma,\nu\,,\ y = s/2m_W^2\,,\ \beta = (1-2/y)^{1/2}\,,$$

$$d_\mathrm{Z} = (s - m_\mathrm{Z}^2 + im\Gamma_\mathrm{Z})^{-1}\,,\ x = \bar{s}_W^2\,,\ \mathrm{e_Z} = \mathrm{e}\sqrt{\frac{1-x}{x}}\,,$$

$$\left.\begin{matrix} g_V^i \\ g_A^i \end{matrix}\right\} = -\frac{\mathrm{e}}{\sqrt{\frac{4x}{1-x}}}\left\{\begin{matrix} T_3^i - 2Q_i x \\ -T_3^i \end{matrix}\right\}\,,$$

$$\begin{aligned}
M_{\mathrm{ZZ}} &= (2x e_\mathrm{Z})^2[(g_V^i)^2 + (g_A^i)^2]s^2|d_\mathrm{Z}|^2 D_1(k_\mathrm{Z}, k_\mathrm{Z}; \lambda_\mathrm{Z}, \lambda_\mathrm{Z})/e^4\,,\\
M_{\gamma\gamma} &= (2xQ_i)^2 D_1(k_\gamma, k_\gamma; \lambda_\mathrm{Z}, \lambda_\mathrm{Z})\,,\\
M_{\gamma\mathrm{Z}} &= 4xQ_i e_\mathrm{Z} g_V^i s\mathrm{Re}\,(d_\mathrm{Z}) D_1(k_\gamma, k_\mathrm{Z}; \lambda_\gamma\lambda_\mathrm{Z})/e^2\,,\\
M_{\mathrm{Z}\nu} &= x e_\mathrm{Z}(g_V^i + g_A^i)s\mathrm{Re}(d_\mathrm{Z}) D_2(k_\mathrm{Z}, \lambda_\mathrm{Z})/e^2\,,\\
M_{\gamma\nu} &= xQ_i D_2(k_\gamma, \lambda_\gamma)\,,\\
M_{\nu\nu} &= D_3\,,
\end{aligned}$$

$$\begin{aligned}
D_1(a,b;c,d) &= 2\beta^2\frac{3z_0 - z_0^3}{3}\{(ab+2cd)y^2\\
&\quad - [1-d-c+(a+c)(b+d)]y + 3\} + 4z_0\beta^2(1+a+c)(1+b+d)y\,,\\
D_2(a,b) &= 4\beta^2\frac{3z_0 - z_0^3}{3}(ay^2 - y) + 4(1+\beta^2)z_0 + 16z_0(1+a+b)\\
&\quad \times(1+\beta^2 y) + 4Lm_W^2\frac{2(1+a+b)+\frac{1}{y}}{y\beta}\,,\\
D_3 &= \beta^2 y^2\frac{3z_0 - z_0^3}{3} + 8z_0 y + 2L\beta^{-9}(1+\beta^2) + \frac{8z_0(1-2\beta^2 z_0^2 + \beta^4)}{\beta^2[(1+\beta^2) - 4\beta^2 z_0^2]}\,,
\end{aligned}$$

$$L = \ln\frac{(1+\beta+2\beta z_0)}{(1+\beta^2 - 2\beta z_0)}\,,$$

z_0 is the detector acceptance (for the ideal detector the z_0 is equal to zero).

The reader should be reminded that the analytical expressions for the cross sections of the reaction (3.8) have been obtained in the works [23–25]. The authors of those works also used the L_{WWV} in the form (3.1). Their results do not coincide with each other. Moreover, none of the obtained expressions for the s-channel cross section does not coincide with that of the paper [26] while mine does at $z_0 = 0$.

The asymptotic behaviour of the σ_T of the process (3.8) is

$$\sigma_\mathrm{T} \simeq s \qquad (3.12)$$

in the CM's. While in the SM and in the EGM's we have

$$\sigma_T \simeq s^{-1} \ln s \,. \tag{3.13}$$

It should be noted that the CM's could give the σ_T with the asymptotic behaviour which is just the same as in the SM. It is possible in the case when the MM's have the structure very similar to that of hadrons. I can make the simplest possible ansatz that all form factors are

$$s^{-n} c_i \theta(\sqrt{s} - \Lambda) + q_i \,, \tag{3.14}$$

where q_i are the MM's values in the SM, c_i are arbitrary constants,

$$\theta(x) = \begin{cases} 0\,, & x > 0 \\ 1\,, & x < 0 \end{cases}$$

$n \geq 1$ for the λ_V and $n \geq 2$ for the k_V ($V = \gamma, Z$) [27].

Further on I shall assume that the MM's are the constant quantities and work at the energies scale $\sqrt{s} < \Lambda$. Then the $(\sigma_T)_{CM}$ asymptotic behaviour of the reaction (3.8) (hereafter I consider the case $f_i = e^-$) will be defined by (3.12). One would think that in these circumstances one will be able to distinguish the CM's from the other models. In fact, $(\sigma_T)_{SM}$ having reached its maximum at $\sqrt{s} \simeq 200$ GeV starts to decrease according to (3.13). Analogously, (I take the case when the parameters values are defined by (2.16) and z_0 is equal to 1) the $(\sigma_T)_{SLR}$ reaches its first maximum at $\sqrt{s} \simeq 194$ GeV, then decreases till $\sqrt{s} \simeq 742$ GeV and afterwards starts to increase till its second maximum. Having reached this maximum at $\sqrt{s} = 800$ GeV the $(\sigma_T)_{SLR}$ decreases (see (3.13)). In the CM's we are expecting that the σ_T will grow as the linear function of s over the whole physical region values of s. However, under certain values of the MM's the $(\sigma_T)_{CM}$ exposes an interesting property easily visible from its analytical expression. The σ_T increases until energy 200 GeV or so, then it falls down up to its minimum and only after that it starts to grow linearly on s [28].

Giving the SM values to the three MM parameters and varying only a single one I can write four simplest conditions for the existence of minima

$$\begin{aligned} -0.78 &< \Delta k_\gamma < 1.10 \\ -0.61 &< \Delta\lambda_\gamma < 0.77 \\ -0.70 &< \Delta k_Z < 0.83 \\ -0.57 &< \Delta\lambda_Z < 0.59 \,. \end{aligned} \tag{3.15}$$

No processes have been observed experimentally in which the MM's occur at the tree level because present colliders do not have sufficient energy or luminosity. However, these MM's occur in loops, and contribute to observed processes. Since in calculation some one-loop diagrams with ultraviolet divergences are involved, one has to fix a procedure to obtain finite results. Then the more preferable scheme is the use of the regulator, because one can consider the parameter Λ which enters into a regulator as the compositeness scale. In paper [29] the neutrino–nucleon scattering, the polarized electron–deuterium asymmetry and the process

$$e^-e^+ \longrightarrow \mu^-\mu^+ \tag{3.16}$$

have been examined at low energies to extract the constraints on the MM's. At $\Lambda = 1$ TeV the results are close to the minima existence conditions of the σ_T of the reaction (1.1), except the inequality for the Δk_Z (the Δk_Z bounds obtained in the paper [29] are $-0.77 < \Delta k_Z < 0$). The coincidence of those results with my own are due to the fact that the investigated cross sections have also the minima when the MM's values are bounded by the obtained constraints. These minima are situated enough far away from the region of the investigated energies. That leads to the small deviations of the CM cross sections from the SM ones

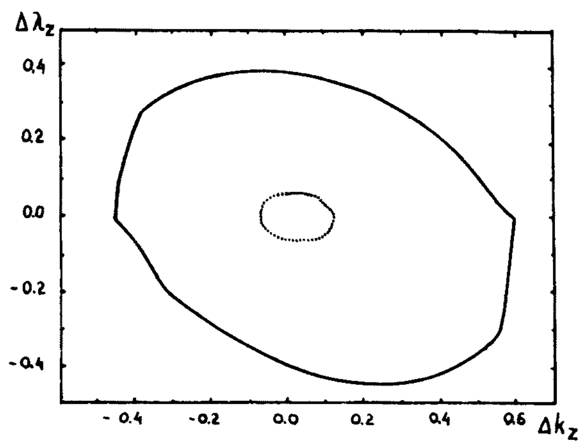

Fig. 7. The bounds on the Δk_Z and the $\Delta\lambda_Z$ following from measurements of the $\sigma_{T(e^-e^+\to W^-W^+)}$ fulfilled with a precision of about 5% at energies up to 200 (solid line) and 500 GeV (dotted line).

In Fig. 7 I present the constraints on the value of the Δk_Z and the $\Delta\lambda_Z$ which follow from the measurements fulfilled with a precision of about5% ($\Delta k_\gamma = \lambda_\gamma = 0$). The values of Δk_Z and $\Delta\lambda_Z$ lie inside the area restricted by the solid line and dashed one for the energies up to 200 and 500 GeV, respectively.

Except of the region of the negative values of Δk_Z my results agree quantitatively with the ones [29] which were also obtained at the same conditions. I am not able to explain the lack of agreement for negative Δk_Z because the authors of the paper [29] have not used any analytical expression for the $\sigma_{T(e^-e^+\to W^-W^+)}$ (the σ_T has been calculated numerically) and have imposed the approximate relation $\Delta\lambda_Z = \Delta\lambda_\gamma$, which has not been assumed by me.

I shall again introduce the quantity δ which is determined according to

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

and compare the δ behaviour for the CM's and the LR models. For this purpose I choose the values of the energy and the SM parameters just the same as in the case of Fig. 2. Again, for the sake of simplicity, I shall give the SM values to the three MM parameters and vary only a single one. Thus, the notation $\delta(\Delta k_Z)$ means that δ is a function of Δk_Z at $\Delta k_\gamma = \Delta\lambda_Z = \Delta\lambda_\gamma = 0$ and so on.

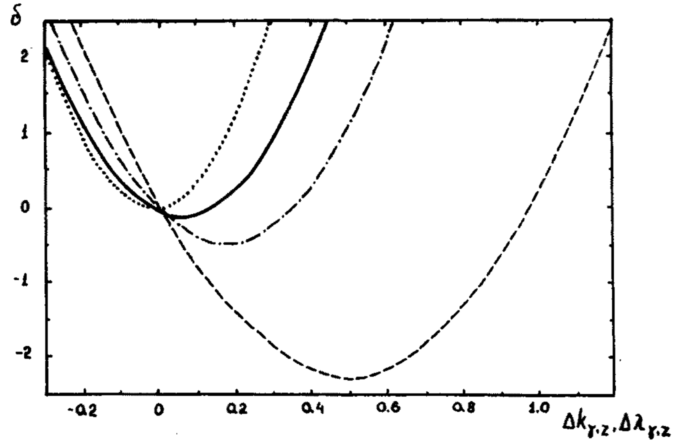

Fig. 8. The δ for the CM as a function of MM's parameters. The solid (dashed) line stands for $\delta(\Delta k_Z)$ ($\delta(\Delta k_\gamma)$) and the dotted (dash–dotted) one does for $\delta(\Delta\lambda_Z)$ ($\delta(\Delta\lambda_\gamma)$).

In Fig. 8 I present the graph of the functions $\delta(\Delta k_\gamma)$, $\delta(\Delta k_\gamma)$, $\delta(\Delta\lambda_Z)$, and $\delta(\Delta\lambda_\gamma)$. We see that the δ is proportional to the deviations of the MM's from their SM values. From Fig. 8 and Fig. 2 also follows that the both models could lead to the same predictions.

For illustrative purposes in Fig. 9 I present the δ dependence on $\sqrt{s}$ at $\Delta k_Z = 1$, $\Delta k_\gamma = \Delta\lambda_\gamma = \Delta\lambda_Z = 0$. In the energies region in which the $(\sigma_T)_{CM}$ is less than the $(\sigma_T)_{SM}$ the maximum deviations of the CM cross

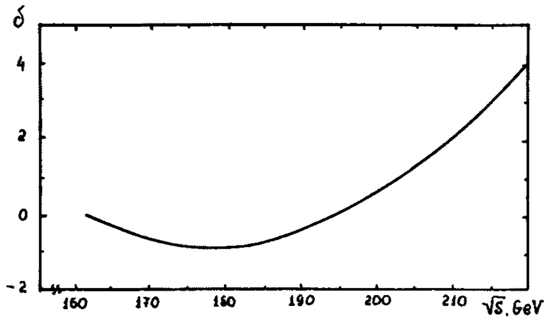

Fig. 9. The δ *versus* $\sqrt{s}$ at $\Delta k_\gamma = 1$ and $\Delta k_Z = \Delta\lambda_Z = \Delta\lambda_\gamma = 0$.

section from the SM one can reach the values of about 0.5%. The inclusion of the RC will influence the degree of these deviations. The oblique RC which are determined by all light fermions are common for the SM and the CM because both the models have some kind of unification. For the CM this kind of unification is provided by the Vector Boson Dominance mechanism and the requirement of a global SU(2) weak isospin symmetry. The other RC can have the origin both similar to one of the SM (virtual exchanges of the ordinary heavy particles, Bremsstrahlung, ...) and different from that of the SM (a shift in the W mass due to the MM's). The latter could give either the increase or the decrease of the CM cross section depending on the sign of the MM parameters. I am going to discuss in detail their role elsewhere.

4. Conclusions

The model which unifies all the possible symmetric and asymmetric LR models has been investigated. At the definite parameters values it reproduces the SM in the sector of the ordinary SM particles. The differential and the total cross sections of the W-pair production in e^-e^+ colliding beams have been calculated. The results of the analysis of the total cross section at LEP II energies could be formulated as follows. The investigation of the cross section in the case of the right-polarized electrons and the left-polarized positrons σ_{lr} will give the answer to the question whether the LR model is the true model of the electroweak interactions or not. If it is then with the help of the unpolarized cross section we can determine the quantity $\Delta\rho_M$, reliably. Then reinvestigating σ_{lr} we could obtain more strict bounds on the mass of the heavy right-handed neutrino, the mixings angles ξ and Φ, and the coupling constant g_R than we had before.

A single production of Z_2-boson in pp collisions has been investigated. It has been shown that this reaction is a good test for the g_R definition.

I have also obtained the total cross sections of (1.1) for the CM's. It has been found that at the definite MM's values the obtained total cross section has the minimum. Its value and its position significantly depends on the MM's values. The comparison with the results of the LR model has shown that the predictions of both the models could coincide into some interval of energies. The measurement of the total cross section of the process (1.1) fulfilled with a precision of about 5% gives us the opportunity to constrain the values of Δk_Z and $\Delta\lambda_Z$ in the interval 28% (12%) and 16% (7%) at the energy up to 200 (500) GeV, respectively.

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REFERENCES

[1] L.F. Abbott, E. Farni, *Phys. Lett.* **B101**, 69 (1981); *Nucl. Phys.* **B189**, 547 (1981).
[2] B. Schrempp, F. Schrempp, *Nucl. Phys.* **B231**, 109 (1984).
[3] E. Ma, *Phys. Rev. Lett.* **63**, 1042 (1989).
[4] R.N. Mohapatra, J.C. Pati, *Phys. Rev.* **D11**, 566 (1975).
[5] R.N. Mohapatra, D. Sidhu, *Phys. Rev. Lett.* **38**, 667 (1977).
[6] R.N. Mohapatra, G. Senjanovic, *Phys. Rev.* **D23**, 165 (1981).
[7] R.M. Francis *et al.*, *Phys. Rev.* **D43**, 2369 (1991).
[8] S. Rajpoot, *Phys. Rev.* **D40**, 3795 (1989).
[9] O. Boyarkin, *Acta Phys. Pol.* **B23**, 1031 (1992).
[10] O. Boyarkin, *Yad. Fiz.* **57**, 135 (1993).
[11] R.N. Mohapatra, *Prog. Part. Nucl. Phys.* **26**, 1 (1991).
[12] J.F. Gunion *et al.*, *Phys. Rev.* **D40**, 546 (1989).
[13] G. Altarelli *et al.*, *Nucl. Phys.* **B342**, 15 (1990); W. Beenakker, W. Hollik, *Z. Phys.* **C40**, 141 (1988).
[14] D.A. Ross, M. Veltman, *Nucl. Phys.* **B95**, 135 (1975).
[15] O. Boyarkin, *Yad. Fiz.* **54**, 839 (1991).
[16] J. Polak, M. Zrałek, *Nucl. Phys.* **B369**, 385 (1991).
[17] G. Altarelli *et al.*, *Phys. Lett.* **B263**, 459 (1991).
[18] E. Eichten, I. Hinchlife, K. Lane, C. Quiggs, *Rev. Mod. Phys.* **56**, 579 (1984); *Rev. Mod. Phys.* **58**, 1065 (1986).
[19] K.J. Gaemers, G.J. Gounaris, *Z. Phys.* **C1**, 259 (1979).
[20] J.C. Wallet, *Z. Phys.* **C30**, 575 (1986).
[21] O. Karnig Mikaelian, *Phys. Rev.* **D30**, 1115 (1984).
[22] M.A. Samuel, J.H. Reid, *Progr. Theor. Phys.* **76**, 184 (1986).
[23] K. Gaemers, M. Velzen, *Z. Phys.* **C43**, 103 (1989).
[24] R. Robinett, *Phys. Rev.* **D28**, 1192 (1983).

[25] A.J. Ramalho, *Phys. Rev.* **D38**, 3574 (1988).
[26] N. Cabibbo, R. Gatto, *Phys. Rev.* **124**, 1577 (1961).
[27] O. Boyarkin, in: Proc. of the XIII Warsaw Symp. on Elementary Particle Physics, Kazimierz, Poland, May 28–June 1, 1990, World Scientific, Singapore, p. 225; *Sov. Phys. J.* No 9, 773 (1990).
[28] O. Boyarkin, *Yad. Fiz.* **54**, 1058 (1991).
[29] G.L. Kane *et al. Phys. Rev.* **D39**, 2617 (1989).