Asymmetric left-right model of electroweak interactions

O. M. Boyarkin
Department of Physics, University of Grodno, Ozheshko Street 22, 230023, Grodno, Belarus
(Received 14 October 1993)

The asymmetric left-right model (LRM) is considered. The differential and the total cross sections of the reaction $e^-e^+\to W^-W^+$ are calculated at the level of the improved Born approximation. The investigation of the unpolarized e^-e^+ beams shows that already at the CERN e^+e^- collider LEP II energies we can either establish or limit such parameters of the LRM as $\Delta\rho_M$ and Φ. It is demonstrated that the LRM manifestations are maximum for $e_R^-e_L^+$ beams. The measurements of the $\sigma_{e_R^- e_L^+ \to W^-W^+}$ will yield interesting information about m_{ν_R} and g_R. The process of the Z_2-boson single production in pp collisions is also considered. It is shown that it displays the extreme sensibility to both the g_R and the Φ values.

PACS number(s): 12.60.Cn, 11.30.Hv, 13.10.+q, 13.38.−b

I. INTRODUCTION

All particle-physics phenomena within the range of energies available today give impressive support to the standard model (SM) of the electroweak interaction based on the $SU(2)_L\times U(1)_Y$ gauge group. Despite its enormous success and its internal consistency and elegance it is widely believed that it is not the ultimate truth. Many other models have been proposed to extend the SM. One of the most attractive extensions of the SM is the left-right model (LRM) based on the $G_{LR}=SU(2)_L\times SU(2)_R\times U(1)_{B-L}$ gauge group. This model accounts for many, but by no means all, physics problems which cannot get a 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 interactions may be interpreted in LRM's as arising out of the spontaneous breaking of parity and the consequent nonvanishing neutrino masses which are possibly required by astrophysics and cosmology. There are also the following important reasons for considering these models: (a) LRM's give a comprehensive picture of the fermion spectrum [1] (from under 20 eV for the electron neutrino to over 100 GeV for the yet undiscovered top quark); (b) the quantum numbers of the group U(1) are identified with $B-L$ (instead of Y, having no physical meaning) which allows us to link the breaking of parity and the breaking of $B-L$; (c) LRM's 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 [2]. All the versions of the LRM's predict the existence of such new particles as W_2 and Z_2 bosons and the massive right-handed neutrino ν_R. So, the search for the left-right (LR) symmetry manifestations could use two different approaches. The former is based on the study of the possible indirect effects of these particles while the latter is based on the search for their direct production at the existing and future planned colliders.

In this paper I consider both approaches. As an example illustrating the former I investigate the reaction

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

for the case of the unpolarized and the polarized e^-e^+ beams. Within the symmetric LRM the unpolarized cross sections were considered in Refs. [3,4]. However, the right-handed neutrino mass m_{ν_R} has not been taken into account in these works. At first glance this assumption seems to be quite natural. In the case of the heavy neutrino,

$$m_{\nu_R}\sim\sqrt{s}\ ,$$

where s is the energy in the center of mass system, we could expect that due to the decoupling theorem (DT) [5] the influence of m_{ν_R} would be negligibly small. Recall that the DT reads, "all the effects connected with the heavy virtual particle, when its mass tends to infinity, are unobservable." It is well known that in the theories with exact gauge symmetry, such as QCD and QED, the DT works without any exceptions. But in the theories with spontaneously broken symmetry the DT could lead to the wrong answer. The examples are the influence of the top quark mass m_t both on B_d^0-$\bar{B}_d^0$ oscillations and the value of $\Delta\rho$. In these cases at $m_t\to\infty$ the effects caused by the t quark are increasing. It turns out that the total cross section of (1.1) also gives us the example of DT violation for the chiral fermions. In order to illustrate the other approach I consider the process

$$p\overset{(-)}{p}\to Z_2+\text{anything}\ , \tag{1.2}$$

using the spirit of the parton models.

The plan of the paper is as follows. In Sec. II I propose the model which unifies the wide class of different versions of the LRM's. Then the total cross sections of reactions (1.1) and (1.2) are presented and the effects of LR symmetry are evaluated. Finally, I summarize my work in Sec. III.

II. LR MODEL

There are a lot of papers in which LRM's are considered [6–10]. For my analysis it is convenient to unify all the LRM versions into one common model. To realize this program I start with the Lagrangian of the form

$$L=-\tfrac{1}{4}(W^{\alpha}_{L\mu\nu}W^{\alpha}_{L\mu\nu}+W^{\alpha}_{R\mu\nu}W^{\alpha}_{R\mu\nu}+B_{\mu\nu}B_{\mu\nu})+i\sum_f(\overline{\Psi}^{(f)}_L D_\mu\gamma_\mu\Psi^{(f)}_L+\overline{\Psi}^{(f)}_R D_\mu\gamma_\mu\Psi^{(f)}_R)+\sum_i|D_\mu\varphi_i|^2+L_Y-V\,, \tag{2.1}$$

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

$$V=\sum_i\lambda_i V^{(2)}(\varphi_i,\varphi_i)+\sum_{i,j,k,l}\mu_{ijkl}V^{(4)}(\varphi_i,\varphi_j,\varphi_k,\varphi_l)\,, \tag{2.2}$$

where λ_i and μ_{ijkl} are constants, and $V^{(2)}(\varphi_i,\varphi_i)$ and $V^{(4)}(\varphi_i,\varphi_j,\varphi_k,\varphi_l)$ are the quadratic and the biquadratic on φ_i terms [their obvious form is defined by symmetries imposed on (2.1)]. I introduce an arbitrary number of the following Higgs multiplets [their numbers T_L, T_R, $(B-L)/2$ are given in parentheses]: (a) doublets $X_L(\frac{1}{2},0,-\frac{1}{2})$, $X_R(0,\frac{1}{2},\frac{1}{2})$; (b) triplets $\delta_L(1,0,1)$, $\delta_R(0,1,1)$; (c) bidoublets $\Phi(\frac{1}{2},\frac{1}{2},0)$. I should like to remind the reader that in the LRM's where fermions pick up their masses through radiative corrections (RC's) one should use several multiplets with the same values of $T_{L,R}$ to prohibit μ-e transitions [11]. Next I suppose that the minimum of the potential V corresponds to the following choice of Higgs field vacuum expectation values (VEV's):

$$\langle X_{L,R}\rangle=\begin{bmatrix}0\\ \nu_{L,R}\end{bmatrix},\quad \langle\delta_{L,R}\rangle=\begin{bmatrix}0\\0\\ \Delta_{L,R}\end{bmatrix},\quad \langle\Phi\rangle=\begin{bmatrix}k & 0\\ 0 & k'\exp(i\omega)\end{bmatrix}. \tag{2.3}$$

$\langle\Phi\rangle$ makes W_L and W_R mix with a CP-violating phase ω. The photon A and the bosons Z_1 and Z_2 in the weak eigenstate basis are related to the neutral gauge bosons W_{3L}, W_{3R}, and B in the G_{LR} basis according to

$$\begin{bmatrix}Z_{1\mu}\\ Z_{2\mu}\\ A_\mu\end{bmatrix}=M\begin{bmatrix}W^L_{3\mu}\\ W^R_{3\mu}\\ B_\mu\end{bmatrix} \tag{2.4}$$

where $M=U_N\Lambda$,

$$U=\begin{bmatrix}\cos\Phi & \sin\Phi & 0\\ -\sin\Phi & \cos\Phi & 0\\ 0 & 0 & 1\end{bmatrix},$$

Φ is the mixing angle of the neutral gauge bosons,

$$\Lambda=\begin{bmatrix}eb(g'^{-2}c_\varphi+a_-g_R^{-1}) & eb(g'^{-2}s_\varphi-a_-g_L^{-1}) & -ebg'^{-1}a_+\\ -bg'^{-1}s_\varphi & bg'^{-1}c_\varphi & -ba_-\\ eg_L^{-1} & eg_R^{-1} & eg'^{-1}\end{bmatrix}$$

$$a_+=g_R^{-1}s_\varphi+g_L^{-1}c_\varphi,\quad a_-=g_R^{-1}c_\varphi-g_L^{-1}s_\varphi,\quad b^{-1}=\sqrt{g'^{-2}+a_-^2}\,,\quad e^{-2}=g_L^{-2}+g_R^{-2}+g'^{-2}\,.$$

So, the neutral-current interaction in this model depends on five parameters which can be taken as VEV's of the Higgs fields, gauge couplings g' and g_R, or as two Z masses, mixing angles Φ and φ, and the Wainberg angle θ_W:

$$\alpha,\beta_L,\beta_R,g_R,g'\rightarrow m_{Z_1},m_{Z_2},\theta_W,\Phi,\varphi\,, \tag{2.5}$$

where

$$\alpha=\tfrac{1}{2}\sum(|k|^2+|k'|^2),\beta_{L,R}=\tfrac{1}{2}\sum(|v_{L,R}|^2+4|\Delta_{L,R}|^2)$$

In the symmetric case the number of independent parameters is reduced to four and the angle φ is their function. This circumstance allows us to use for φ the term "the LR symmetry-violating angle." Changing φ I can reproduce all the known LRM's. For example, the symmetric LRM's proposed in Refs. [6] and [7] are reproduced at $\varphi=-\pi/4$ and $\varphi=0$, respectively, and the asymmetric LRM of Ref. [10] is reproduced at $\varphi=\arctan[g_L g_R(g_R^{-2}+g'^{-2})]$. It is important to note here that if the role of φ was only reduced to the redefinition of the mixing angle Φ then we would take, for example, any LR model (say, the model of Ref. [7] with matrix Λ_0) and obtain all versions of LRM's with arbitrary matrix Λ from it by changing Φ. It would be true if the condition

$$U_N\Lambda=U'_N\Lambda_0 \tag{2.6}$$

took place, where U'_N has the form

$$U'_N=\begin{bmatrix}\cos\Phi' & \sin\Phi' & 0\\ -\sin\Phi' & \cos\Phi' & 0\\ 0 & 0 & 1\end{bmatrix}.$$

Elementary algebra implies that condition (2.6) is not satisfied. The angle ϕ is connected with the orientation of the $SU(2)_R$ generator in the group space and it redistributes the role between Φ and m_{Z_2}. For a different angle

φ we have different experimental bounds on Φ and m_{Z_2}. Only the experiment can answer which φ value is the true one. For example, in the case $g_L = g_R$ the analysis of Z_1 decay parameters being made at the CERN e^+e^- collider LEP I gives the following bound on φ [12]:

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

The interaction Lagrangian has the following form in the model under consideration:

$$L_{WWV} = i\rho_{kl}^{(V)}(W_{k\mu\nu}^{*}V_{\nu} + W_{k\nu}V_{\nu\mu}^{*})W_{l\mu} , \tag{2.7}$$

$$L_{\rm NC} = i\sum_f \overline{\Psi}_f \gamma_\mu \left[qA_\mu + \frac{1}{2}\sum_{n=1}^{2}(g_{Vn}^f + g_{An}^f\gamma_5)Z_{n\mu}\right]\Psi_f , \tag{2.8}$$

$$L_{CC} = \frac{ig_L}{2\sqrt{2}} \sum_{\substack{j=e,d \\ i=\nu,u}} (K_{ij})_L \overline{\Psi}_j \gamma_\mu (1+\gamma_5)\Psi_i W_\mu^L + (L \leftrightarrow R) , \tag{2.9}$$

where $V_{\mu\nu} = \partial_\mu V_\nu - \partial_\nu V_\mu \quad (V_\mu = A_\mu, Z_{k\mu}), \quad k = 1,2,$ $W_\mu^L = W_{1\mu}\cos\xi - W_{2\mu}\exp(-i\omega)\sin\xi,$

$W_\mu^R = W_{1\mu}\exp(i\omega)\sin\xi + W_{2\mu}\cos\xi,\ \psi_{\nu_L} = \nu\cos\alpha - N\sin\alpha,$

$\psi_{\nu_R} = N\cos\alpha + \nu\sin\alpha ,$

$$\rho_{kl}^{(A)} = q\delta_{kl} , \tag{2.10}$$

$$\rho_{ll}^{(Z_1)} = \cos^2\left[\xi + \frac{\pi}{2}\delta_{l2}\right]g_L M_{11} + \sin^2\left[\xi + \frac{\pi}{2}\delta_{l2}\right]g_R M_{12} , \tag{2.11}$$

$$\rho_{kl}^{(Z_1)} = \rho_{lk}^{(Z_1)} = \tfrac{1}{2}\sin 2\xi(g_L M_{11} - g_R M_{12}), \quad k \neq l , \tag{2.12}$$

$$g_{Vk}^f = \left[T_{3L}^f g_L M_{k1} + T_{3R}^f g_R M_{k2} + \left[1 - \frac{\delta_{f\nu}}{2}\right](B-L)g' M_{k3}\right] , \tag{2.13}$$

$$g_{Ak}^f = [T_{3L}^f g_L M_{k1} - T_{3R}^f g_R M_{k2} - \delta_{f\nu} g'(T_{3L}^f - T_{3R}^f)M_{k3}] , \tag{2.14}$$

$(K_{ij})_{L,R}$ are the elements of the Kobayashi-Maskawa matrix, q is the charge of the particle, ξ is the mixing angle of gauge charged bosons, ν and N are Majorana or Dirac spinors describing the mass eigenstates, α is the mixing angle of the light and the heavy neutrinos in the vacuum, and the expressions for $\rho_{ll}^{(Z_2)}$, $\rho_{kl}^{(Z_2)}$, and $\rho_{lk}^{(Z_2)}$ follow from (2.11) and (2.12) by the substitutions

$$g_L M_{11} \to g_R M_{22}, \quad g_R M_{12} \to g_L M_{21} ,$$
$$\xi \to \xi + \pi/2$$

respectively.

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

$$\varphi = \Phi = \xi = 0, \quad g_L = g_R = es_W^{-1}, \quad g' = e(c_W^2 - s_W^2)^{-1/2} \tag{2.15}$$

in the symmetric case and

$$\varphi = \Phi = \xi = 0, \quad g_L = es_W^{-1},$$
$$g' = \pm(c_W^2 e^{-2} - g_R^{-2})^{-1/2}, \quad |g_R| \rangle g_L s_W c_W^{-1} \tag{2.16}$$

in the asymmetric one. In conclusion, in this case, I should like to pay some attention to the conditions imposed on the coupling constants (CC's) g_L, g_R, and g'. If we start to build our theory from the G_{LR} gauge group directly, then the CC's 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 CC's all must come together at the grand unification scale, the value of which should be consistent with the proton decay constraint. These demands impose limitations on the choice of the unifying group. For example, the nonsupersymmetric SU($N \geq 5$) models are ruled out while the supersymmetric SU(5) and SO($N \geq 10$) models satisfy the demands above. The Q^2 dependence of the CC's can be calculated from the renormalization group equations in which these quantities enter quadrically. The choice of both the GUT and the scale of the underlying symmetry breaking defines the conditions imposed on the CC's at the electroweak scale (see, for example, the general LR model [13] in which the CC's are defined by

$$g_L = es_W^{-1}, \quad g' = (c_W^2 e^{-2} - g_R^{-2})^{-1/2}, \quad g_L^2/2 \leq g_R^2 \leq g_L^2$$

and from the very beginning the LR symmetry-violating angle was taken to be equal to zero). Up to now we cannot confidently give preference to the definite GUT. Therefore, investigating the G_{LR} gauge theory as a low energy approximation of the GUT we should consider the most general version of the LRM. As we see, the proposed theory satisfies this demand. It possesses a larger space in fitting its parameters to the experimental data. However, further on I shall, for the sake of simplicity, ignore the possibilities connected with the variation of the LR symmetry violation angle and take it to be equal to zero. I shall also suppose that CC's are defined by condition (2.16).

As I am going to take into account the effects connected with the right-handed neutrino I would like to make some remarks about its properties. The value of the mass of the right-handed neutrino 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 m_{ν_R} could be obtained from cosmological considerations as well as weak decay processes. For example, if the heavy neutrinos are Majorana particles the known bound on the neutrinoless double β decay half-life of ^{76}Ge gives the lower bound [14]

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

The upper bound obtained from the arguments of vacuum stability [13] practically does not depend on m_{W_R} and it is about 1 TeV.

The basic decay modes of the right-handed neutrino are

$$\nu_R \to W_i^{\pm} + 1^{\mp} , \tag{2.17}$$

$$\nu_R \to \nu_L + V , \tag{2.18}$$

where $V = \gamma, Z_{1,2}$ and $i = 1,2$. Their decay widths are defined as

$$\Gamma_{\nu_R \to W_i + l} \cong \frac{g_R^2 h_i m_{\nu_R}}{32\pi} \left[y_i + \frac{2}{y_i^2} - \frac{3}{y_i} \right] , \tag{2.19}$$

where $y_i = (m_{\nu_R}/m_{W_i})^2$,

$$h_i = \begin{cases} \sin^2\xi, & i=1, \\ \cos^2\xi, & i=2 , \end{cases}$$

for (2.17), and

$$\Gamma_{\nu_R \to \nu_L + V} \cong \frac{m_{\nu_R}^5}{2\pi} \left[\sum_{i=1}^{2} f_i \right] \tau_V , \tag{2.20}$$

where

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

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

and m_l is the mass of the conjugated lepton, for (2.18) [the process (2.18) with $V=\gamma$ was considered in Ref. [15]].

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 (2.17) with $i=1$ is very much suppressed by the factor $\sin^2\xi$. So, for the case $m_{W_2} > m_{\nu_R}$ which I shall consider, in what follows, shall be able to neglect $\Gamma_{\nu_R \to \text{all}}$ for the sake of simplicity.

In the LRM at $e^- e^+$ annihilation we may observe the processes

$$e^- e^+ \to W_k^- W_n^+ . \tag{2.21}$$

In Fig. 1 the Feynman diagrams contributing to the process (2.21) are shown. I shall be limited by including the RC at the level of the improved Born approximation [16]. I remind that the obtained results will be the same both for the Dirac and Majorana neutrinos. Then the differential cross section for the polarized initial and the unpolarized final particles is defined by

$$\begin{aligned} \frac{d\sigma^{(kn)}}{dt} = \frac{\rho}{(4\pi s)^2} \Bigg\{ \Bigg[& \frac{1-\lambda\bar\lambda}{4} \left[|e^2[1+(-1)^{k+n}] - s\sum_{l=1}^{2} \rho_{kn}^{(Z_l)} g_{Vl}^e d_{Z_l}|^2 + |2s \sum_{l=1}^{2} \rho_{kn}^{(Z_l)} g_{Al}^e d_{Z_l}|^2 \right] \\ & - \frac{\lambda - \bar\lambda}{4} \left[2se^2[1+(-1)^{k+n}] \sum_{l=1}^{2} \rho_{kn}^{(Z_l)} g_{Al}^e \operatorname{Re}(d_{Z_l}) - 2s^2 \sum_{l=1}^{2} \rho_{kn}^{(Z_l)} g_{Al}^e d_{Z_l} \sum_{j=1}^{2} \rho_{kn}^{(Z_j)} g_{Vl}^e d_{Z_j}^* \right] \Bigg] B_1^{(kn)}(s,u,t) \\ & - \Bigg[\frac{(1-\bar\lambda)(1+\lambda)}{4} g_L^2 a_+^{kn} \left[e\rho_{kn}^{(A)} - \frac{s}{2} \sum_{l=1}^{2} \rho_{kn}^{(Z_l)} (g_{Vl}^e + g_{Al}^e) \operatorname{Re}(d_{Z_l}) \right] \\ & - \frac{(1+\bar\lambda)(1-\lambda)}{4} g_R^2 a_-^{kn} \left[e\rho_{kn}^{(A)} - \frac{s}{2} \sum_{l=1}^{2} \rho_{kn}^{(Z_l)} (g_{Vl}^e - g_{Al}^e) \operatorname{Re}(d_{Z_l}) \right] d_{\nu_R} \Bigg] B_2^{(kn)}(s,t,u) \\ & + \left[\frac{(1+\bar\lambda)(1-\lambda)}{8} g_R^4 b_-^{kn} d_{\nu_R}^2 + \frac{(1-\bar\lambda)(1+\lambda)}{8} g_L^4 b_+^{kn} \right] B_3^{kn}(s,t,u) \Bigg\} \end{aligned} \tag{2.22}$$

where

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

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

$$q = e$$

$d_{\nu_R} = t/(t - m_{\nu_R}^2)$, $\lambda(\bar\lambda)$ is the helicity of the electron (positron), s, u, and t are the ordinary Mandelstam variables, and the functions $B_i^{(kn)}(s,u,t)$ $(i=1,2,3)$ are defined in Ref. [4].

The quantity ρ entering (2.22) is caused by the RC's which are defined by all heavy particles and the ρ value is determined by

$$\rho = 1 + \Delta\rho_t + \Delta\rho_M + \cdots , \tag{2.23}$$

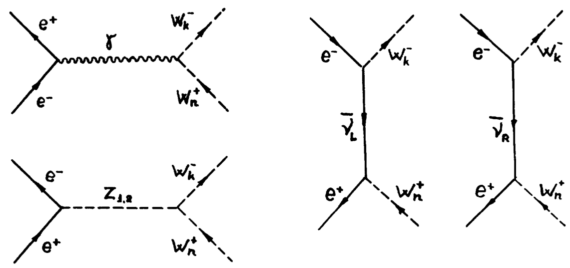

FIG. 1. Diagrams corresponding to the processes $e^-e^+ \rightarrow W_k^- W_n^+$.

where $\Delta\rho_t \simeq 3G_F m_t^2/8\sqrt{2}\pi$. $\Delta\rho_M$ arises due to the mixing in the gauge vector boson sector and has the general asymptotic form [16]

$$\Delta\rho_M = c_0^2(m_{Z_1}/m_{Z_2})^2 - c_1^2(m_{W_1}/m_{W_2})^2 \tag{2.24}$$

with c_0 and c_1 constants depending on the Higgs particles VEV's and the CC's. Recall that when $\Delta\rho_M$ does not equal zero then the effective s_W^2 in the LRM is connected with the $\bar{s}_W^2 = \sin^2\theta_W$ of the SM by the relation

$$s_W^2 = \bar{s}_W^2 - [\bar{s}_W^2 \bar{c}_W^2/(\bar{c}_W^2 - s_W^2)]\Delta\rho_M \ .$$

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 boson mass and for $m_H \gg m_{W_1}$ it has a negative sign. The contribution connected with the presence of weak isotriplet Higgs bosons is [17]

$$(\Delta\rho_{\rm th})_{L,R} = -\frac{2\Delta_{L,R}^2}{v^2}$$

where $v = 246$ GeV is the standard Higgs-doublet VEV. I shall, in what follows, make the simplifying assumption that all these additional contributions to ρ are mutually canceled. The constraint on ρ coming from the measurements of m_Z at LEP and m_W/m_Z at UA2 and the Collider Detector at Fermilab (CDF) collaborations has the form

$$\rho = 1.0066 \pm 0.0058 \ .$$

So, the upper and lower bounds on $\Delta\rho_M$ coming from the experiments are

$$(\Delta\rho_M)_{\rm upper, lower} = (\Delta\rho)_{\rm upper, lower} - \Delta\rho_t - 1 \ . \tag{2.25}$$

From (2.22) it follows that all the effects connected with Z_2 and ν_R are negligible when

$$m_{Z_2} \rightarrow \infty \ , \quad m_{\nu_R} \rightarrow \infty \ .$$

So, in this case the DT works very well. However, for the total cross section the situation with ν_R is drastically altered.

The expression for the total cross section $\sigma^{(kn)}$ follows from (2.22) after multiplication by $\pi\beta s$ and the replacement

$$\begin{aligned} &B_i^{(kn)}(s,u,t) \rightarrow D_i^{(kn)}, \quad i = 1,2,3, \\ &B_j^{(kn)}(s,u,t)(d_{\nu_R})^{j-1} \rightarrow D_j^{(kn)}, j = 2,3, \end{aligned} \tag{2.26}$$

where

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

$$\begin{aligned} D_{2\nu}^{(kn)} = &\frac{(s - m_n^2 - m_k^2)}{m_n^2 m_k^2} \left\{ 2(m_n^2 + m_k^2) + \frac{s\beta^2}{6} + \frac{2m_n^2 m_k^2}{s} - \frac{m_{\nu_R}^2}{2s}\left[2m_{\nu_R}^2 - m_n^2 - m_k^2 + s + \frac{8m_n^2 m_k^2}{m_n^2 + m_k^2 - s} \right] \right\} \\ &+ \frac{4L}{\beta s}\left[-m_n^2 - m_k^2 - \frac{m_n^2 m_k^2}{s} + m_{\nu_R}^2(m_n^2 + m_k^2 - s)\left[\frac{m_n^2 + m_k^2}{2m_n^2 m_k^2} + \frac{3}{4s} - \frac{m_{\nu_R}^4}{4sm_n^2 m_k^2} \right] + \frac{\beta^2 s m_{\nu_R}^4}{4m_n^2 m_k^2} \right], \end{aligned}$$

$$\begin{aligned} D_{3\nu}^{(kn)} = &\frac{1}{12m_n^2 m_k^2}\left[s^2\beta^2 - 48m_n^2 m_k^2 + 12s(m_n^2 + m_k^2) - 6m_{\nu_R}^2(s - m_k^2 - m_n^2 + 4m_{\nu_R}^2) \right] \\ &+ C_\nu \frac{s m_{\nu_R}^4 (m_n^2 + m_k^2)}{m_n^2 m_k^2} \\ &+ \frac{L}{\beta s}\left\{ 2(s - m_n^2 - m_k^2) + 5m_{\nu_R}^2 + \frac{m_{\nu_R}^2}{2m_n^2 m_k^2}\left[m_{\nu_R}^2(4m_{\nu_R}^2 - 3m_n^2 - 3m_k^2 + 3s) - 4s(m_n^2 + m_k^2) \right] \right\}, \end{aligned}$$

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

$$L = \ln\left| \frac{-m_n^2 - m_k^2 + 2m_{\nu_R}^2 + \beta s + s}{-m_n^2 - m_k^2 + 2m_{\nu_R}^2 - \beta s + s} \right| \ .$$

As we see from (2.26) in $\sigma^{(kn)}$ there are the addenda which are proportional to $m_{\nu_R}^{2q}$ $(q=1,2)$. They arise when we integrate in (2.22) the expressions containing $(td_{\nu_R})^q$. In the case of unpolarized e^-e^+ beams, their contribution to $\sigma^{(11)}$ is very small because the dominant addenda caused by ν_L exchange enter $\sigma^{(11)}$ as $D_2^{(11)}\cos^2\xi$ and $D_3^{(11)}\cos^4\xi$ while the ones determined by ν_R exchange are $D_{2\nu}^{(11)}\sin^2\xi$ and $D_{3\nu}^{(11)}\sin^4\xi$.

For the comparison with the SM it is convenient to introduce the quantity δ characterizing the experimental sensitivity. I define it by

$$\delta^{\lambda\bar{\lambda}}(x)=\frac{\Delta\sigma}{\sqrt{(\sigma)_{\rm SM}}}\sqrt{LT}\ , \tag{2.27}$$

where LT is the integrated luminosity of the collider in units of pb^{-1}, $\Delta\sigma=(\sigma)_{\rm LRM}-(\sigma)_{\rm SM}$, and $(\sigma)_{\rm SM(LRM)}$ is the total cross section in the SM (LRM). The $\delta^{\lambda\bar{\lambda}}(x)$ is an observable of the effect from new physics and it gives the deviations from the SM expressed in standard error units. Previous work on this problem [4] has shown that in the LEP II energy region the total cross section of the LRM calculated at the tree level differs from that of the SM on 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 $\Delta\rho_M$. Recall that because of the structure of $\Delta\rho_M$ given by Eq. (2.24), $\Delta\rho_M$ could have not only positive and the negative values but zero ones also. In order to receive an idea about the values of deviations caused by the RC I shall use the upper, lower, and zero bounds of $\Delta\rho_M$ in my analysis of the obtained total cross section.

First I consider the case of the symmetric LRM. In Fig. 2 I represent the $\delta(\sqrt{s}\,)$ (hereafter the absence of the superscripts λ and $\bar{\lambda}$ means the case of the unpolarized e^-e^+ beams) of reaction (1.1) for $(\Delta\rho_M)_{\rm upper}$, $(\Delta\rho_M)_{\rm lower}$, and $(\Delta\rho_M)=0$ at $LT=500\ \mathrm{pb}^{-1}$. I note that at the given values of m_t $(\Delta\rho_M)_{\rm lower}$ is negative. In numerical calculations I used the following values of the structural parameters (SP's) of the SM and the LRM [18,19]:

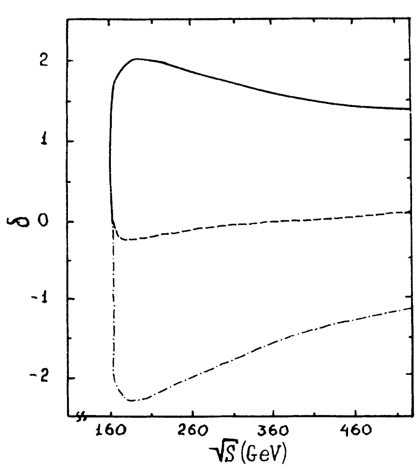

FIG. 2. δ for the symmetric LRM as a function of $\sqrt{s}$ for $(\Delta\rho_M)_{\rm upper}$ (solid line), $(\Delta\rho_M)_{\rm lower}$ (dash-dotted line), and $\Delta\rho_M=0$ (dashed line).

$$m_{Z_1}=80.13\ \mathrm{GeV},\quad m_{W_1}=90.177\ \mathrm{GeV},\quad \bar{s}_W=0.23,$$
$$m_t=145\ \mathrm{GeV},\quad m_{z_2}=800\ \mathrm{GeV},\quad m_{W_2}=477\ \mathrm{GeV},$$
$$m_{\nu_R}=400\ \mathrm{GeV},\quad \Phi=9.6\times10^{-3},\quad \xi=3.1\times10^{-2}\ . \tag{2.28}$$

We see that the possible deviations of the LRM from the SM lie within the region restricted by the curves $(\Delta\rho_M)_{\rm upper}$ and $(\Delta\rho_M)_{\rm lower}$. Next I shall demand that δ be greater than the 2σ uncertainty on the total cross section, which means that the LR symmetry manifestations will be at 95% the confidence level (C.L.) As follows from Fig. 2 only in the energy region $\sqrt{s}\le 260$ GeV may one hope to observe the clean signal. On further increasing the energy, $\delta(\sqrt{s}\,)$ decreases and then starts to grow near Z_2 resonance. At $\Delta\rho_M>0$ and the chosen value Φ [according to today's estimates [19] the bounds on Φ are -0.005 $(-0.0048)<\Phi<0.0099(0.0096)$ at $m_H=100$ GeV (1 TeV) and $m_t=90$ GeV in both cases] $\delta(\sqrt{s}\,)$ has a minimum in the energy region less than m_{Z_2}. The presence or absence of this minimum is defined by the Φ value only. At the bigger values of $\Phi(\Phi\ge4\times10^{-2})$ the minimum would be absent and $\delta(\sqrt{s}\,)$ would monotonically increase up to $\sqrt{s}=m_{Z_2}$. The analysis shows that the minimum position $(\sqrt{s}\,)_{\rm min}$ and its value $\delta_{\rm min}(\sqrt{s}\,)$ significantly depend on the chosen values m_{Z_2} and Φ; in contrast, the change of the other SP's weakly influences both $\delta_{\rm min}(\sqrt{s}\,)$ and $(\sqrt{s}\,)_{\rm min}$. At $\Delta\rho_M\le0$ the function $\delta(\sqrt{s}\,)$ crosses the axis of the abscissas at 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 energy region is the fast increases of its modulus with the growth of $|\Delta\rho_M|$. Therefore, 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. The LR symmetry effects become more significant on elimination of the diagram with the ν_L exchange. They are maximum in the case of completely right-polarized electrons $(\lambda=-1)$ and left-polarized positrons $(\bar{\lambda}=1)$ (RL polarization). For example, when $\Delta\rho_M=0$ and $m_{\nu_R}=100$ GeV the deviation of the LRM from the SM reaches the 2σ level for $\delta_+^{-1,0}(\sqrt{s}\,)$ [the subscript $+$ $(-)$ will mean that I consider the positive (negative) g' values in (2.16)] and $\delta_+^{-1,1}(\sqrt{s}\,)$ at $\sqrt{s}\cong420$ and $\sqrt{s}\cong590$ GeV, respectively. Nothwithstanding the fact that the RL polarized cross section σ^{RL} is much less than the unpolarized one σ (for example, at LEP II energy, $\sigma^{RL}/\sigma\cong10^{-2}$), the σ^{RL} investigation would play a decisive role for the existence of the LRM. In Fig. 3 I represent $\delta^{-1,1}(\sqrt{s}\,)$ for the different values of m_{ν_R} and $\Delta\rho_M$ and two signs of g' (hereafter the other SP's of the LRM and LT are the same as in the case of Fig. 2). It should be noted here that $|\delta_+^{\lambda\bar{\lambda}}(\sqrt{s}\,)|=|\delta_-^{\lambda\bar{\lambda}}(\sqrt{s}\,)|$. The $\delta^{-1,1}(\sqrt{s}\,)$ is also sensitive to the values of the mixing angles ξ and Φ. However, just the same as for $\delta(\sqrt{s}\,)$, the dependence of

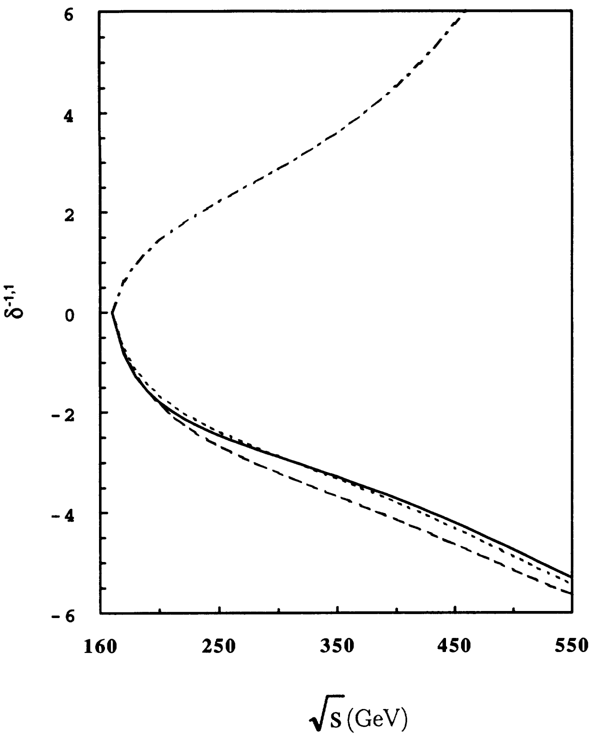

FIG. 3. $\delta_{\mp}^{-1,1}$ versus $\sqrt{s}$ for the cases (a) $m_{\nu_R}=100$ GeV and $(\Delta\rho_M)_{\text{upper}}$ (solid line), (b) $m_{\nu_R}=400$ GeV and $(\Delta\rho_M)_{\text{upper}}$ (dashed line), and (c) $m_{\nu_R}=100$ GeV and $(\Delta\rho_M)_{\text{lower}}$ (dotted line). The function $\delta_{-}^{-1,1}(\sqrt{s})$ with $(\Delta\rho_M)_{\text{lower}}$ and $m_{\nu_R}=400$ GeV is represented by the dash-dotted line.

$\delta^{-1,1}(\sqrt{s})$ on Φ is much stronger than the dependence on ξ. For example, at $\Delta\rho_M=0$ the variation of Φ from 9.6×10^{-3} up to 9.6×10^{-4} yields a decrease of $\delta_{\mp}^{-1,1}(\sqrt{s})$ of one order of magnitude, while the variation of ξ from 3.1×10^{-2} up to 3.1×10^{-3} yields an increase of $\delta_{\mp}^{-1,1}(\sqrt{s})$ of a few percent. It is immediately apparent that in the asymmetric case $\delta^{-1,1}(\sqrt{s})$ could be used for g_R definition. Figure 4 shows the behavior of $\delta_{\mp}^{-1,1}(g_R/g_L)$ for different values of m_{ν_R} and $\Delta\rho_M$ at $\sqrt{s}=196$ and $\sqrt{s}=500$ GeV. The results for LEP II are very disappointing because the deviations are smaller than the statistical errors except in a region of very specific values of g_R/g_L. The situation is much better for ILC.

Having the total cross section I can define the longitudinal polarization asymmetry as

$$A_{LR}=\frac{(\sigma)_{1,0}-(\sigma)_{-1,0}}{(\sigma)_{1,0}+(\sigma)_{-1,0}}\ , \tag{2.29}$$

where $(\sigma)_{\lambda,\bar{\lambda}}$ is the total cross section either of the LRM or the SM in the case when the electron and the positron have the helicities λ and $\bar{\lambda}$, respectively. Recall that this quantity is sensitive to the pure weak RC and can be measured with a precision of $1/\sqrt{N_W P_{e^-}}$ [N_W is the number of W bosons producing in the reaction (1.1) and P_{e^-} is the degree of longitudinal polarization of the electron]. In Fig. 5 I display A_{LR} as a function of $\sqrt{s}$ for the case of the symmetric LRM. For comparison the SM re-

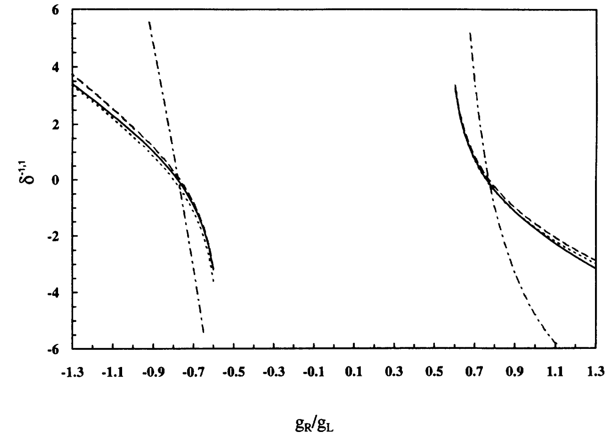

FIG. 4. The $\delta_{\mp}^{-1,1}$ as a function of g_R/g_L at $\sqrt{s}=196$ GeV for the cases (a) $(\Delta\rho_M)_{\text{upper}}$ and $m_{\nu_R}=100$ GeV (dashed line); (b) $(\Delta\rho_M)_{\text{lower}}$ and $m_{\nu_R}=100$ GeV (dotted line); and (c) $(\Delta\rho_M)_{\text{upper}}$ and $m_{\nu_R}=400$ GeV (solid line). At $\sqrt{s}=500$ GeV, $\Delta\rho_M=0$, and $m_{\nu_R}=100$ GeV the function $\delta_{\mp}^{-1,1}(g_R/g_L)$ is represented by the dash-dotted line.

sults are also represented. Again we see that no clear signal will be seen at LEP II energy. The situation does not improve for such machines as ILC and the Next Linear Collider (NLC), since with the growth of the energy the cross section decreases and as a result the statistical errors increase. The analysis shows that the quantity A_{LR} is practically insensitive to the variations of g_R and m_{ν_R}.

Hadron colliders provide a nice opportunity for the investigation of extended gauge models. For example, the extra neutral gauge boson discovery limits at the CERN Large Hadron Collider (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 g_R definition. It appears that the single production of a Z_2 boson is a good tool for this purpose. I adopt the spirit of the parton model. Then in the lowest order of the Drell-Yan approximation the total cross section of the process

$$ab\rightarrow Z_n+\text{anything}\ , \tag{2.30}$$

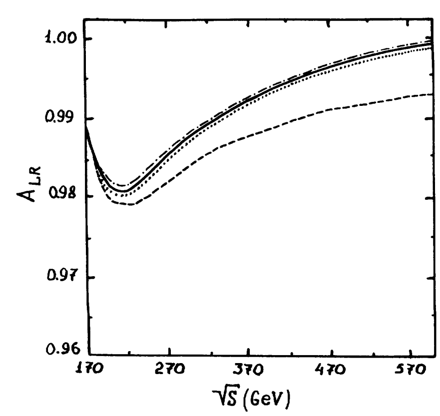

FIG. 5. A_{LR} as a function of $\sqrt{s}$ for $(\Delta\rho_M)_{\text{upper}}$ (dash-dotted line), for $(\Delta\rho_M)_{\text{lower}}$ (dotted line), and for $(\Delta\rho_M)=0$ (solid line). The results of the SM are represented by the dashed line.

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

$$\sigma=\frac{1}{24m_{Z_2}^2\sqrt{1+\beta^2}}\sum_i\left[(g_{V2}^{q_i})^2+(g_{A2}^{q_i})^2\right]\tau\frac{dL_{q_iq_i}}{d\tau}\ , \tag{2.31}$$

where $\tau=m_{Z_n}^2/s$, $s=(p_a+p_b)^2$, $f_q^{(a)}(x,Q^2)$ is the distribution function of the quark flavor i in hadron a, the parameter Q^2 of which the 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 $\tau dL_{q_i\bar{q}_i}/d\tau$ is defined by the expression

$$\tau\frac{dL_{q_i\bar{q}_i}}{d\tau}=\int_\tau^1\left[f_{q_i}^{(a)}(x,Q^2)f_{\bar{q}_i}^{(b)}\left[\frac{\tau}{x},Q^2\right]+(q_i\leftrightarrow\bar{q}_i)\right]\frac{dx}{x}\ . \tag{2.32}$$

From (2.31) and the expressions for $g_{V2}^{q_i}$ and $g_{A2}^{q_i}$ it follows that the cross section of (2.30) is practically insensitive to the choice of the g_R sign.

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

$$\sigma=\frac{(\sigma B)_{\rm ALRM}-(\sigma B)_{\rm SLRM}}{\sqrt{(\sigma B)_{\rm SLRM}}}\sqrt{LT} \tag{2.33}$$

where B is the Z_2 branching ratio and the subscripts ALRM (SLRM) mean that the quantity in parentheses is calculated within the asymmetric (symmetric) LRM. At present δ gives the deviations from the symmetric LRM expressed in standard error units. In order to eliminate backgrounds I shall only consider the following leptonic modes of Z_2 decay:

$$Z_2\rightarrow e^-e^+,\mu^-\mu^+\ . \tag{2.34}$$

In Fig. 6 is shown the δ of the reaction (2.30) versus g_R/g_L for LHC energies ($\sqrt{s}=17$ TeV , $LT=10^4$ pb^{-1}) in the case of pp collisions. 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. [20] (set 2). I also used the following values of parameters:

$$m_{Z_2}=600\ \text{GeV},\quad \phi=9.6\times10^{-3}\ .$$

Under these conditions, the typical event sample for the LHC will consist of $N_{Z_2}\cong2\times10^6$. This is the main reason which accounts for the extreme sensitivity of δ to the g_R variation. It is obvious that the reaction (2.30) is also a good test for the Φ definition.

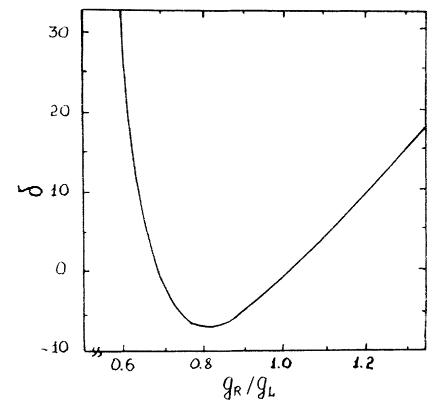

FIG. 6. δ versus g_R/g_L for pp collisions at the LHC.

III. CONCLUSIONS

A model which unifies all the possible symmetric and asymmetric LRM's has been investigated. At definite parameter values it reproduces the SM in the sector of ordinary SM particles. The analysis of W_1-pair production in e^-e^+ colliding beams has been carried out on the basis of exact calculations for the total cross section. For unpolarized e^-e^+ beams the LR symmetry manifestations could be observable in the energy region up to $\sqrt{s}\cong260$ GeV only. In this region, depending on the values of $\Delta\rho_M$ and Φ, the deviations of the LRM from the SM could reach the 2σ level. At the near-vanishing values of $\Delta\rho_M$ the only way of observing the indirect signals of the LR symmetry is the measurement of the cross section at a polarized-beam machine with energy $\sqrt{s}\geq400$ GeV. The investigation has shown that the indirect effects should be more significant for $e_R^-e_L^+$ beams. In this case the values of the deviations from the SM predictions depend mainly on Φ and $\sqrt{s}$. The effects connected with the heavy right-handed neutrino are maximum for the RL polarization. Unfortunately their values are too small to be measured reliably. The role of ν_R becomes essential for the processes

$$e^-e^+\rightarrow W_1^-W_2^+,W_2^-W_2^+ \tag{3.1}$$

by virtue of the fact that the $W_2\nu_Re$ coupling is proportional to $\cos\xi$. Thus the observation of the processes (3.1) would prove to be a powerful tool for probing the properties of ν_R.

However, if nature realized the LRM with the angle Φ of the order of 10^{-3} and less, then the LR symmetry would only display itself in the direct production of the new particles predicted by the LRM. I have considered a single production of a Z_2 boson in pp collisions. It has been shown that this reaction is extremely sensitive to the variations of g_R and Φ.

ACKNOWLEDGMENTS

I would like to thank F. Feruglio for the fruitful discussion of the DT role in the theories with spontaneously broken symmetry.

[1] Erneset Ma, Phys. Rev. Lett. **63**, 1042 (1989).
[2] R. N. Mohapatra and J. C. Pati, Phys. Rev. D **11**, 566 (1975).
[3] O. Boyarkin, Acta Phys. Polon. B **21**, 495 (1990); O. Boyarkin, Yad. Fiz. **54**, 839 (1991) [Sov. J. Nucl. Phys. **54**, 506 (1991)].
[4] O. Boyarkin, Sov. J. Nucl. Phys. **54**, 506 (1991).
[5] T. Appelquist and C. Bernard, Phys. Rev. D **22**, 200 (1980); A. C. Longhitano, *ibid.* **22**, 1166 (1980).
[6] R. N. Mohapatra and D. Sidhu, Phys. Rev. Lett. **38**, 667 (1977).
[7] R. N. Mohapatra and G. Senjanovic, Phys. Rev. D **23**, 165 (1981).
[8] R. M. Francis *et al.*, Phys. Rev. D **43**, 2369 (1991).
[9] S. Rajpoot, Phys. Rev. D **40**, 3795 (1989).
[10] M. Cvetic and B. W. Lynn, Phys. Rev. D **35**, 51 (1987).
[11] Gwo-Guang Wong, Phys. Rev. D **46**, 3994 (1992).
[12] O. Boyarkin, Acta Phys. Polon. B **23**, 1031 (1992).
[13] R. N. Mohapatra, Prog. Part. Nucl. Phys. **26**, 1 (1991).
[14] R. N. Mohapatra, Phys. Rev. D **34**, 909 (1986).
[15] R. N. Mohapatra and S. Nussinov, Phys. Rev. D **39**, 1378 (1989).
[16] G. Altarelli *et al.*, Nucl. Phys. **B342**, 15 (1990); W. Beenakker and W. Hollik, Z. Phys. C **40**, 141 (1988).
[17] D. A. Ross and M. Veltman, Nucl. Phys. **B95**, 135 (1975).
[18] J. Polak and M. Zralek, Nucl. Phys. **B369**, 385 (1991).
[19] G. Altarelli *et al.*, Phys. Lett. B **263**, 459 (1991).
[20] E. Eichten *et al.*, Rev. Mod. Phys. **56**, 579 (1984); **58**, 1065 (1986).