Measurement of the ν_τ magnetic moment at CERN LEP in a left-right symmetric model

A. Gutiérrez-Rodríguez
Escuela de Física, Universidad Autónoma de Zacatecas
Apartado postal C-580, 98068 Zacatecas, Zac., Mexico

M.A. Hernández and A. Rosado
Instituto de Física, Universidad Autónoma de Puebla
Apartado postal J-48, Col. San Manuel, 72570 Puebla, Pue., Mexico

M. Maya
Facultad de Ciencias Físico Matemáticas, Universidad Autónoma de Puebla
Apartado postal 1364, 72000 Puebla, Pue., Mexico

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We perform a theoretical calculation of the ν_τ magnetic moment through the reaction $e^+e^- \rightarrow \nu\,\bar{\nu}\,\gamma$, in the framework of a left-right symmetric model at LEP energies. We find that the bound is almost independent of the mixing angle ϕ of the model in the allowed experimental range for this parameter.

Keywords: Neutrino electromagnetic properties

Se realiza un cálculo teórico del momento magnético del ν_τ a través de la reacción $e^+e^- \rightarrow \nu\,\bar{\nu}\,\gamma$, en el marco de un modelo con simetría izquierda-derecha a energías alcanzables en LEP. Se encuentra una cota que es casi independiente del ángulo de mezcla ϕ del modelo en el intervalo de valores experimentales permitido para este parámetro.

Descriptores: Propiedades electromagnéticas del neutrino

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1. Introduction

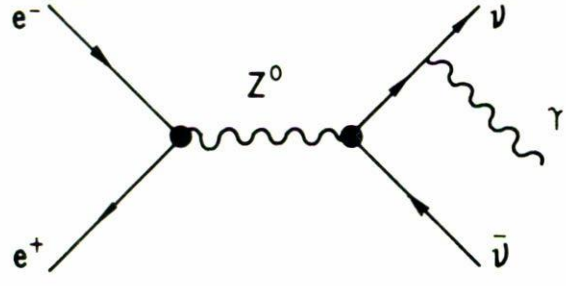
Neutrinos seem to be likely candidates for carrying features of physics beyond the Standard Model (SM) [1]. Apart from masses and mixings also magnetic moments and electric dipole moments are signs of new physics and are of relevance in terrestrial experiments, the solar neutrino problem, astrophysics and cosmology [2]. At the present time, all the available experimental data for electroweak processes can be well understood in the context of the SM, except the results of the KAMIOKANDE experiment on the neutrino mass [3]. Hence, the SM is the starting point of all the extended gauge models. In other words, any gauge group with physical sense must have as a subgroup the $\mathrm{SU}(2)_L \times \mathrm{U}(1)$ group of the standard model, in such way that their predictions agree with those of the SM at low energies. The purpose of the extended theories is to explain some fundamental aspects which are not clarified in the frame of the SM. One of these aspects is the origin of the parity violation at the current energies. The Left-Right Symmetric Models (LRSM) based on the $\mathrm{SU}(2)_R \times \mathrm{SU}(2)_L \times \mathrm{U}(1)$ gauge group [4] give an answer to that problem, since restore the parity symmetry at high energies and give their violations at low energies as a result of the breaking of gauge symmetry. Detailed discussions on LRSM can be found in the literature [5–7].

In 1994, T.M. Gould and I.Z. Rothstein [8] reported a bound on the tau neutrino magnetic moment, which they obtained trough the analysis of the process $e^+e^- \rightarrow \nu\bar{\nu}\gamma$, near the Z^0-resonance, by considering a massive tau neutrino and using Standar Model Ze^+e^- and $Z\nu\bar{\nu}$ couplings.

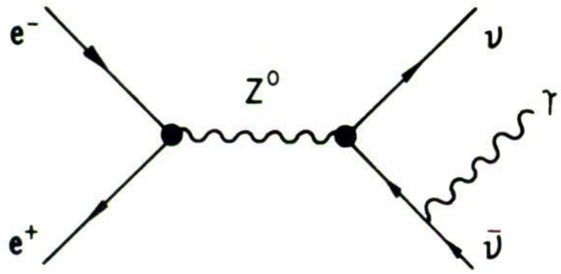
At low center of mass energy $s \ll M^2_{Z^0}$, the dominant contribution to the process $e^+e^- \rightarrow \nu\bar{\nu}\gamma$ involves the exchange of a virtual photon [9]. The dependence on the magnetic moment comes from a direct coupling to the virtual photon, and the observed photon is a result of initial state Bremsstrahlung.

At higher s, near the Z^0 pole $s \approx M^2_{Z^0}$, the dominant contribution for $E_\gamma > 10$ GeV [10] involves the exchange of a Z^0 boson. The dependence on the magnetic moment now comes from the radiation of the observed photon by the neutrino or antineutrino in the final state. The Feynman diagrams which give the most important contribution to the cross section are shown in Fig. 1. We emphasize here the importance of the final state radiation near the Z^0 pole, which occurs preferentially at high E_γ compared to conventional Bremsstrahlung.

Our aim in this paper is to analize the reaction $e^+e^- \rightarrow \nu\bar{\nu}\gamma$ with recent data from LEP [10–12] near the Z^0 boson resonance in the framework of a left-right symmetric model and attributing a magnetic moment to a massive tau neutrino.



(a)



(b)

FIGURE 1. The Feynman diagrams contributing to the process $e^+e^- \to \nu\bar{\nu}\gamma$, in left-rihgt symmetric models.

Processes measured near the resonance have served to bound the tau neutrino magnetic moment. In this paper we take advantage of this fact to set bounds for $\kappa(\nu_\tau)$ for different values of the mixing angle ϕ [13–15], which is in agreement with other constraints previously reported [8, 9].

We will do our analysis near the resonance of the Z^0 ($s \approx M^2_{Z^0}$). As a consequence our results are independent of the mass of the additional heavy Z^0_R gauge boson which appears in this kind of models and so we have the mixing angle ϕ between the left and the right bosons as the only additional parameter, besides the SM parameters.

This paper is organized as follows. In Sect. 2 we describe the model with the Higgs sector having two doublets and one bidoublet. In Sect. 3 we perform the calculation of the matrix elements and the differential cross section of the process $e^+e^- \to \nu\bar{\nu}\gamma$. In Sect. 4 we present and discuss our numerical results for the total cross section as a function of the mixing angle ϕ. Finally, a summary of our conclusions is given in Sect. 5.

2. The left-right symmetric model (LRSM)

We consider a left-right symmetric model (LRSM) with one bidoublet Φ and two doublets χ_L, χ_R whose vacuum expectation values break the gauge symmetry to give mass to the left and right heavy gauge bosons. This is the origin of the parity violation at low energies [5], that is, at energies available at present accelerators. The Lagrangian for the Higgs sector of the LRSM is [6]

$$\mathcal{L}_{\mathrm{LRSM}} = (D_\mu \chi_L)^\dagger (D^\mu \chi_L + (D_\mu \chi_R)^\dagger (D^\mu \chi_R) + Tr(D_\mu \Phi)^\dagger (D^\mu \Phi). \quad (1)$$

The covariant derivates are written as

$$\begin{aligned} D_\mu \chi_L &= \partial_\mu \chi_L - \frac{1}{2} i g \tau \cdot \mathbf{W}_L \chi_L - \frac{1}{2} i g' B \chi_L, \\ D_\mu \chi_R &= \partial_\mu \chi_R - \frac{1}{2} i g \tau \cdot \mathbf{W}_R \chi_R - \frac{1}{2} i g' B \chi_R, \\ D_\mu \Phi &= \partial_\mu \Phi - \frac{1}{2} i g (\tau \cdot \mathbf{W}_L \Phi - \Phi \tau \cdot \mathbf{W}_R). \end{aligned} \quad (2)$$

In this model there are seven gauge bosons: the charged $W^1_{L,R}$, $W^2_{L,R}$ and the neutral $W^3_{L,R}$, B. The coupling constants for the left and right sector are equal $g_L = g_R$, since manifest left-right symmetry is assumed [16].

The transformation properties of the Higgs bosons under the group $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)$ are $\chi_L \sim (1/2, 0, 1)$, $\chi_R \sim (0, 1/2, 1)$ and $\Phi \sim (1/2, 1/2^*, 0)$. After spontaneous symmetry breaking, the Higgs bosons develop a vacuum expectation value of the form

$$\begin{aligned} \langle \chi_L \rangle &= \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_L \end{pmatrix}, \\ \langle \chi_R \rangle &= \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_R \end{pmatrix}, \\ \langle \Phi \rangle &= \frac{1}{\sqrt{2}} \begin{pmatrix} k & 0 \\ 0 & k \end{pmatrix}, \end{aligned} \quad (3)$$

breaking the symmetry group to the $U(1)_{em}$ and giving mass to the gauge bosons and fermions. Only the photon remains massless. The part of the Lagrangian that contains the mass terms for the charged bosons is

$$\mathcal{L}^C_{\mathrm{mass}} = (W^+_L \;\; W^+_R) M^C \begin{pmatrix} W^-_L \\ W^-_R \end{pmatrix}, \quad (4)$$

where $W^\pm = \frac{1}{\sqrt{2}}(W^1 \mp W^2)$.

The mass matrix M^C is

$$M^C = \frac{g^2}{4} \begin{pmatrix} v_L^2 + k^2 + k'^2 & -2kk' \\ -2kk' & v_R^2 + k^2 + k'^2 \end{pmatrix}, \quad (5)$$

this matrix is diagonalized by an orthogonal transformation which is parametrized [16] by an angle ζ. This angle has been restricted to have a very small value from the hyperon β decay data [17].

Similarly, the part of the Lagrangian that contains the mass terms for the neutral bosons is

$$\mathcal{L}^N_{\mathrm{mass}} = \frac{1}{8} (W^3_L \, W^3_R \, B) M^N \begin{pmatrix} W^3_L \\ W^3_R \\ B \end{pmatrix}, \quad (6)$$

where the matrix M^N is given by

$$M^N = \frac{1}{4}\begin{pmatrix} g^2(v_L^2 + k^2 + k'^2) & -g^2(k^2 + k'^2) & -gg' v_L^2 \\ -g^2(k^2 + k'^2) & g^2(v_R^2 + k^2 + k'^2) & -gg' v_R^2 \\ -gg' v_L^2 & -gg' v_R^2 & g'^2(v_L^2 + v_R^2) \end{pmatrix}. \tag{7}$$

Since the process $e^+e^- \to \nu\bar{\nu}\gamma$ is neutral, we center our attention to the mass terms of the Lagrangian for the neutral sector, [Eq. (6)].

The matrix M^N for the neutral gauge bosons is diagonalized by an orthogonal transformation which can be written in terms of the angles θ_W and ϕ [18]

$$U^N = \begin{pmatrix} c_W c_\phi & -s_W t_W c_\phi - r_W s_\phi / c_W & t_W(s_\phi - r_W c_\phi) \\ c_W s_\phi & -s_W t_W s_\phi + r_W c_\phi / c_W & -t_W(c_\phi + r_W s_\phi) \\ s_W & s_W & r_W \end{pmatrix}, \tag{8}$$

where $c_W = \cos\theta_W$, $s_W = \sin\theta_W$, $t_W = \tan\theta_W$ and $r_W = \sqrt{\cos 2\theta_W}$, with θ_W being the electroweak mixing angle. Here, $c_\phi = \cos\phi$ and $s_\phi = \sin\phi$. The angle ϕ can be considered as the angle that mixes the left and right handed neutral gauge bosons $W^3_{L,R}$. The expression that relates the left and right handed neutral gauge bosons $W^3_{L,R}$ and B with the physical bosons Z_1, Z_2 and the photon is

$$\begin{pmatrix} Z_1 \\ Z_2 \\ A \end{pmatrix} = U^N \begin{pmatrix} W_L^3 \\ W_R^3 \\ B \end{pmatrix}. \tag{9}$$

The diagonalization of (5) and (7) gives the mass of the charged $W^\pm_{1,2}$ and neutral $Z_{1,2}$ physical fields:

$$M^2_{W_{1,2}} = \frac{g^2}{8}[v_L^2 + v_R^2 + 2(k^2 + k'^2) \mp \sqrt{(v_R^2 - v_L^2)^2 + 16(kk')^2}], \tag{10}$$

$$M^2_{Z_1,Z_2} = B \mp \sqrt{B^2 - 4C}, \tag{11}$$

respectively, with

$$B = \frac{1}{8}[(g^2 + g'^2)(v_L^2 + v_R^2) + 2g^2(k^2 + k'^2)],$$

$$C = \frac{1}{64}g^2(g^2 + 2g'^2)[v_L^2 v_R^2 + (k^2 + k'^2)(v_L^2 + v_R^2)].$$

Taking into account that $M^2_{W_2} \gg M^2_{W_1}$, from the expressions for the masses of M_{Z_1} and M_{Z_2} we conclude that the relation $M^2_{W_1} = M^2_{Z_1}\cos^2\theta_W$ still holds in this model.

From the Lagrangian of the LRSM we extract the terms for the neutral interaction of a fermion with the gauge bosons $W^3_{L,R}$ and B:

$$\mathcal{L}^N_{\rm int} = g(J^3_L W^3_L + J^3_R W^3_R) + \frac{g'}{2} J_Y B. \tag{12}$$

Explicitly the interaction Lagrangian for $Z_1 \to f\bar{f}$ is [19]

$$\mathcal{L}^N_{\rm int} = \frac{g}{c_W} Z_1 \Big[\Big(c_\phi - \frac{s_W^2}{r_W}s_\phi\Big)J_L - \frac{c_W^2}{r_W}s_\phi J_R\Big], \tag{13}$$

where the left (right) current for the fermions are

$$J_{L,R} = J^3_{L,R} - \sin^2\theta_W J_{\rm em},$$

and

$$J_{\rm em} = J^3_L + J^3_R + \frac{1}{2}J_Y,$$

is the electromagnetic current. From (13) we can find the amplitude $\mathcal{M}$ for the decay of the Z_1 boson with polarization ϵ^λ into an fermion-antifermion par:

$$\mathcal{M} = \frac{g}{c_W}[\bar{u}\gamma^\mu \frac{1}{2}(ag_V - vg_A\gamma_5)v]\epsilon^\lambda_\mu, \tag{14}$$

with

$$a = c_\phi - \frac{s_\phi}{r_W},$$

$$v = c_\phi + r_W s_\phi. \tag{15}$$

In the following section we perform the calculation of the differential cross section for the reaction $e^+e^- \to \nu\bar{\nu}\gamma$ by using the expression (14) for the transition amplitude.

3. The differential cross sections

The expression for the amplitude $\mathcal{M}$ of the process $e^+e^- \to \nu\bar{\nu}\gamma$ due to only Z^0 exchange, according to the diagrams depicted in Fig. 1, and using the expression for the amplitude given in Eq. (14) and assuming that a massive Dirac neutrino is characterized by two phenomenological parameters, *i.e.* a magnetic moment μ_{ν_τ} (expressed in units of the electron Bohr magnetons) and a charge radius $\langle r^2 \rangle$ is given by

$$\mathcal{M}_a = \frac{-g^2}{8\cos^2\theta_W(l^2 - m_\nu^2)} \times [\bar{u}(p_3)\Gamma^\alpha(\not{l} + m_\nu)\gamma^\beta(a - v\gamma_5)v(p_4)] \times \frac{(g_{\alpha\beta} - p_\alpha p_\beta/M_Z^2)}{[(p_1+p_2)^2 - M_Z^2 - i\Gamma_Z^2]} \times [\bar{u}(p_2)\gamma^\alpha(ag_V - vg_A\gamma_5)v(p_1)]\epsilon^\lambda_\alpha, \tag{16}$$

$$\mathcal{M}_b = \frac{-g^2}{8\cos^2\theta_W(k^2 - m_\nu^2)} \times [\bar{u}(p_3)\gamma^\beta(a - v\gamma_5)(\not{k} + m_\nu)\Gamma^\alpha v(p_4)] \times \frac{(g_{\alpha\beta} - p_\alpha p_\beta/M_Z^2)}{[(p_1+p_2)^2 - M_Z^2 - i\Gamma_Z^2]} \times [\bar{u}(p_2)\gamma^\alpha(ag_V - vg_A\gamma_5)v(p_1)]\epsilon_\alpha^\lambda, \quad (17)$$

with

$$\Gamma^\alpha = eF_1(q^2)\gamma^\alpha + \frac{ie}{2m_\nu}F_2(q^2)\sigma^{\alpha\mu}q_\mu, \quad (18)$$

the neutrino electromagnetic vertex, where q is the momentum transfer and $F_{1,2}(q^2)$ dimensionless structure functions. Explicitly

$$F_1(q^2) = \frac{1}{6}q^2\langle r^2\rangle, \quad \text{electroweak charge radius,}$$

$$F_2(q^2) = \mu_\nu\frac{m_\nu}{m_e}, \quad \text{magnetic moment anomalous,}$$

while the coupling constants are given by

$$a = \cos\phi - \frac{\sin\phi}{\sqrt{\cos 2\theta_W}},$$

$$v = \cos\phi + \sqrt{\cos 2\theta_W}\sin\phi, \quad (19)$$

where ϕ is the mixing parameter of the LRSM [13, 14] and ϵ_α^λ is the polarization vector of the photon. l (k) stands by the momentum of the virtual neutrino (antineutrino). Finally, we take $g_V = -\frac{1}{2} + 2\sin^2\theta_W$ and $g_A = -\frac{1}{2}$, according to the experimental data [20].

Using the same notation as in Ref. 8, we find that the magnetic moment coupling as well as the mixing angle ϕ parameter of the LRSM give a contribution to the differential cross section for the process $e^+e^- \to \nu\bar{\nu}\gamma$ of the form

$$\frac{d\sigma}{E_\gamma dE_\gamma d\cos\theta_\gamma} = \frac{\alpha^2\kappa^2}{192\pi}\mu_B^2 C[\phi, x_w] \times F[\phi, s, E_\gamma, \cos\theta_\gamma], \quad (20)$$

where E_γ, $\cos\theta_\gamma$ are the energy and scattering angle of the photon.

The kinematics is contained in the function

$$F[\phi, s, E_\gamma, \cos\theta_\gamma] \equiv \frac{(a^2 + v^2)(s - 2\sqrt{s}E_\gamma) + a^2E_\gamma^2\sin^2\theta_\gamma}{M_{Z^0}^2\Gamma_{Z^0}^2}. \quad (21)$$

The coeficient C is given by

$$C[\phi, x_w] \equiv \frac{a^2/2 + v^2/2 - 4a^2x_w + 8a^2x_w^2}{x_w^2(1-x_w)^2}, \quad (22)$$

with $x_w \equiv \sin^2\theta_W$.

Taking the limit when the mixing angle $\phi = 0$, the expression for a and v is reduced to $a = v = 1$ and the Eq. (20) reduces to the expression (3) given in Ref. 8.

TABLE I

Exp.	$\mathcal{L}(pb^{-1})$	$\theta_{\min}, \theta_{\max}$	$E_{\min}$(GeV)	Events	Ref.
L3 1990	3.8	45, 135	10	0	10
1991	11.2	45, 135	22.8	0	11
ALEPH 1991	8.3	42, 138	17	0	12

TABLE II

ϕ	$\kappa(10^{-6})$
−0.009	5.64
−0.005	5.59
0	5.5
0.004	5.46

TABLE III

ϕ	$\kappa(10^{-6})$
−0.009	3.96
−0.005	3.93
0	3.9
0.004	3.87

4. Results

In order to evaluate the integral of the differential cross section as a function of mixing angle ϕ, in the interval important of each experiment we require cuts on photon angle and energy in order to avoid divergences when the integral is evaluated. We show in Table I, the cuts used in our computations.

According to the experimental data, the allowed range for the mixing angle between Z_L^0 and Z_R^0 is

$$-9\times 10^{-3} \leq \phi \leq 4\times 10^{-3}, \quad (23)$$

with a 90% C.L. [13–15].

As it was discussed in Ref. 8, $N \approx \sigma(\phi,\kappa)\mathcal{L}$ has to be less than 2.2. Using this fact we can put a bound for the tau neutrino magnetic moment as a function of the ϕ mixing parameter. We show the value of this bound for some values of the ϕ parameter in Table II.

These results compare favorably with the bounds obtained in the Refs. 8 and 9. However the derived bounds in Table II could be improved by including data from the entire Z^0 resonance [8] as is shown in Table III.

We end this section plotting the total cross section in Fig. 2 as function of the mixing angle ϕ, for the bounds of the magnetic moment giving in Tables II and III. We observe in these figures that for $\phi = 0$ we reproduce the data previously reported for T.M. Gould and I.Z. Rothstein [8]. Also we observe that the total cross section increases constantly and reaches its maximun value for $\phi = 0.004$.

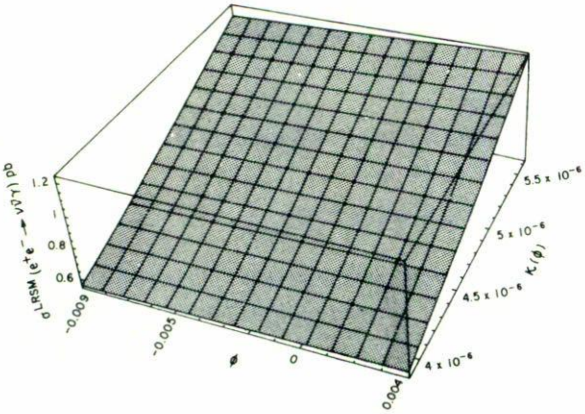


FIGURE 2. The total cross section for $e^+e^- \to \nu\bar{\nu}\gamma$ as function of ϕ $\kappa(\phi)$ (Tables II and III).

5. Conclusions

We have determined a bound on the magnetic moment of a massive tau neutrino in the framework of a left-right symmetric model as a function of the mixing angle ϕ, as is shown in Table II and Table III.

Other upper limits on the tau neutrino magnetic moment reported in the literature are:

$$\mu_{\nu_\tau} < 4.1 \times 10^{-6}\mu_B \quad 90\%\,\text{C.L.}, \tag{24}$$

$$\mu_{\nu_\tau} < 3.6 \times 10^{-6}\mu_B \quad 68\%\,\text{C.L.}, \tag{25}$$

$$\mu_{\nu_\tau} < 5.4 \times 10^{-7}\mu_B \quad 90\%\,\text{C.L.}. \tag{26}$$

The bound given by (24) comes from the search for energetic single photon production in Z^0 decays from L3 collaboration [21]. From LEP-I data on Z^0 partial widths is obtained the bound given by (25). In Ref. 23, Cooper-Sarkar *et al.* have shown how the BEBC (Big European Buble Chamber) beam dump elastic scattering experiments also provide a bound on the tau neutrino diagonal moment of a stable ν_τ given by (26), thus severely restricts the cosmological annihilation scenario [24]. Our results of the Table III confirm those bounds given by (24) and (25).

Our bounds apply to Dirac as well as Majorana transition moments. However, transition moments involving ν_e and ν_μ flavors are more strongly bounded by other accelerator experiments. These limits are $\kappa_{\text{tran}} < 1.08 \times 10^{-10}$ for ν_e, and $\kappa_{\text{tran}} < 7.4 \times 10^{-10}$ for ν_μ [25]. However, a beam dump search for radiative decays gives a limit on the ν_τ transition magnetic moment of $\kappa_{\text{tran}} < 1.1 \times 10^{-9}(\text{MeV/m}_{\nu_\tau})^2$ [26].

In summary, we conclude that the found bound for the tau neutrino magnetic moment is almost independent of the experimental allowed values of the ϕ parameter of the model. In the limit $\phi = 0$ our bound takes the value previously reported by T. M. Gould and I. Z. Rothstein [8].

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