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Cosmological Bounds on a Left-Right Symmetric Model

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ABSTRACT

The right-handed neutrino cross sections of the processes $\nu_L+\ell\leftrightarrow\nu_R+\ell$ and $\nu_R+\ell\leftrightarrow\nu_R+\ell$, where ℓ is a relativistic lepton, are calculated in the model $SU(2)_L\times SU(2)_R\times U(1)$. According to cosmological criterion the parameters of the model are bounded. In particular we obtain the bound $M_{W_R}\geq 30M_{W_L}$, assuming the neutrinos are Dirac four-components particles.

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I. INTRODUCTION

In the standard model of electroweak interactions, the processes involving right-handed neutrinos, ν_R, are proportional to its mass.[1] The importance of the ν_R interactions in the Early Universe becomes apparent when one considers the contribution of weakly interacting particles to ρ, the energy density of the Universe.

There are two kinds of massive neutrinos, Dirac neutrinos (ν_D) and Majorana neutrinos (ν_M). The contribution to the energy density of the Universe is proportional to the number of degrees of freedom. Therefore, we have the relation $\rho_{\nu_D}=2\rho_{\nu_M}$, since the number of degrees of freedom of the Majorana neutrinos are diminished a factor of two by the condition $\nu_M=\nu_M^c$.

In this note we assume the neutrinos are Dirac particles and we study their interaction with relativistic leptons in a left-right symmetric model, $SU(2)_L\times SU(2)_R\times U(1)$.[2]

The contribution to ρ coming from ν_R within the standard model is proportional to m_ν and therefore is negligible. This is the reason we are considering a left-right symmetric model where ν_L and ν_R interact with the same strength. We will see that considering the above left-right symmetric model allows us to constrain the parameters of the model in a more severe form that the low energy analysis of the weak interactions.[3]

In section II we calculate the relevant cross sections involving ν_R and we evaluate them. In section III we give our conclusions.

II. CALCULATION OF σ_{LR} AND σ_{RR}

The Feynman diagrams in Fig. 1 are those that allow the ν_R to be heated by the plasma and remain in equilibrium with the rest of the matter during the lepton epoch.[4] In the left-right symmetric model $SU(2)_L \times SU(2)_R \times U(1)$, the leptons are in doublets $\psi_L \equiv (1/2,0,-1)$, $\psi_R \equiv (0,1/2,-1)$ and the electric charge is given by $Q \equiv t_{3L}+t_{3R}+y/2$.

The neutrinos are four component particles. Both left and right components interact with the gauge bosons with equal strength.

The mass eigenstates are $W_1^{\pm}$, $W_2^{\pm}$, Z_1, Z_2 with masses M_{W_1}, M_{W_2}, M_{Z_1}, M_{Z_2}, respectively. In this case we have double the number of gauge bosons in the standard model. $SU(2)_L \times SU(2)_R \times U(1)$ breaks down to $SU(2)_L \times U(1)_{e.m.}$ in such a way that $M_{W_1} \ll M_{W_2}$ and $M_{Z_1} \ll M_{Z_2}$.[2] In the limit of M_{W_2} and M_{Z_2} going to infinity and zero mixing between left and right currents we recover the standard model. In this limit the processes in Fig. 1 are proportional to the neutrino mass which we are neglecting, hence they must vanish. In this note we are going to neglect the processes in which the ν_R's interact with the leptons via the Higgs bosons. In fact, the Higgs boson-lepton coupling is proportional to the lepton's mass and therefore is negligible in this approximation.

The interaction describing the charged processes in this model can be written as

$$\mathcal{L} = \frac{g}{\sqrt{8}}[J_{L\rho}W_L^{\rho} + J_{R\rho}W_R^{\rho} + h.c.] \tag{1}$$

where $J_{L\rho} = (V-A)_{\rho}$ and $J_{R\rho} = (V+A)_{\rho}$. The mass eigenstates in the charged boson sector are related to W_L and W_R through the rotation

$$\begin{aligned} W_1 &= W_L\cos\zeta - W_R\sin\zeta \\ W_2 &= W_L\sin\zeta + W_R\cos\zeta \end{aligned} \tag{2}$$

where the mixing angle ζ, is proportional to the vacuum expectation values of the Higgs bosons, with the property that it vanishes when the ratio M_{W_1}/M_{W_2} does. (Senjanović, 1979). The converse is not true.

The effective Hamiltonian, at low energies compared with the gauge boson masses, is

$$H_{eff}^{c} = \frac{4G_F}{\sqrt{2}}[aJ_L^{\dagger c}J_L^{\dagger c} + b(J_L^{\dagger c}J_R^{c} + J_R^{\dagger c}J_L^{c}) + cJ_R^{\dagger c}J_R^{c}] \quad , \tag{3}$$

where $a \equiv (\cos^2\zeta + \delta\sin^2\zeta)$, $b \equiv 1/2(\delta-1)\sin^2\zeta$, $c \equiv (\sin^2\zeta + \delta\cos^2\zeta)$, and $\delta \equiv (M_{W_1}/M_{W_2})^2$. Here the left and right currents are written in the form

$$J^{c\alpha}_{L,R} = \bar{\ell}[\gamma^{\alpha}\ \tfrac{1}{2}(1\mp\gamma_5)]\ell \quad .$$

In the limit δ and ζ equal to zero, the (V-A) limit of the standard model is recovered, a=1, b and c are zero and we get

$$H^{c}_{eff} = \frac{G_F}{\sqrt{2}}(V-A)_{\rho}(V-A)^{\rho} \quad .$$

In the neutral current case, we can use the Georgi-Weinberg theorem and express the effective interaction in the following way

$$H^{N}_{eff} = \tfrac{1}{2}[M^{-2}_{LL}(\bar{\ell}\gamma_{\mu}N_{L}\ell)^{2} + M^{-2}_{LR}(\bar{\ell}\gamma_{\mu}N_{L}\ell\bar{\ell}\gamma^{\mu}N_{R}\ell +$$

$$+ \bar{\ell}\gamma_{\mu}N_{R}\ell\bar{\ell}\gamma^{\mu}N_{L}\ell) + M^{-2}_{RR}(\bar{\ell}\gamma_{\mu}N_{R}\ell)^{2}] \quad , \tag{4}$$

where

$$M^{-2}_{LL} \simeq M^{-2}_{W_1} , \qquad M^{-2}_{LR} = M^{-2}_{RR} \simeq M^{-2}_{W_2} \quad ,$$

and

$$N_{L,R} = g[T_3\ \tfrac{1}{2}(1\mp\gamma_5) - \sin^2\theta_W Q] \quad . \tag{5}$$

In the eq. (5), T_3 and Q are the isospin third component and charge of the lepton in consideration, respectively. In this case, we get also the low energy limit of the standard model when M_{W_2} approaches to infinity.

The H^N_{eff} can be expressed in the following way

$$H^N_{eff} = \frac{4G_F}{\sqrt{2}}[J^N_L J^N_L + \delta(J^N_L J^N_R + J^N_R J^N_L)] + \delta(J^N_R J^N_R)] \quad , \tag{6}$$

where

$$J^{N\alpha}_{L,R} = \bar{\ell}\gamma^\alpha[T_3 \frac{1}{2}(1\mp\gamma_5) - \sin^2\theta \; Q]\ell$$

In the eq. (6) we have to notice that the term containing the product $J^N_L J^N_R + J^N_R J^N_L$ is not equal to $2J^N_L J^N_R$ since we must write in fact $J^{\nu N}_L J^{\ell N}_R + J^{\nu N}_R J^{\ell N}_L$. Also we can see in this case there is not a mixing angle explicitly written since the whole contribution coming from the neutral boson masses and the mixing angle can be expressed only in terms of $(M_{W_1}/M_{W_2})^2$ which is δ.

The evaluation of the cross sections from Fig. 1 is straightforward. We have to distinguish between the different processes given by the reactions $\nu_L \ell \leftrightarrow \nu_R + \ell$ (diagram b in Fig. 1) and $\nu_R + \ell \leftrightarrow \nu_R + \ell$ (diagrams a and c in Fig. 1) which give σ_{LR} and σ_{RR}, respectively.

The results are

$$\sigma_{LR} = \frac{\sigma_{LL}}{12}(\delta-1)^2(\sin 2\zeta)^2 \quad , \tag{7}$$

$$\sigma_{RR} = \frac{\sigma_{LL}}{a^2}(c^2 + cd + \frac{1}{3}\, d^2) \quad , \tag{8}$$

where $\sigma_{LL} = 4G_F^2 a^2 E^2/\pi$, $c \equiv (\sin^2\zeta + \delta\cos^2\zeta)$ and $d \equiv \delta(-1/2 + 2\sin^2\theta_W)$.

One can see from eqs. 7 and 8 that both cross sections vanish when one takes the limit $\delta=\zeta=0$. This is because we are working in the limit in which all the leptons under consideration are relativistic and therefore we can neglect the contribution of the lepton mass.

In the phenomenological region of interest, i.e., $\delta\approx\zeta\approx 0$ the cross sections σ_{LR} and σ_{RR} have a similar behavior under $\tan\zeta$ and δ, respectively. In this region we obtain the relations

$$\sigma_{LR}/\sigma_{LL} \simeq \tan^2\zeta \quad , \qquad \sigma_{RR}/\sigma_{LL} \simeq \delta^2 \quad .$$

The ratios σ_{LR}/σ_{LL} and σ_{RR}/σ_{LL} cannot be bigger than 10^{-6} because in this case ν_R will contribute too much to the energy density of the Universe and have modified considerably the present H_e^4 ratio.[5,6]

Assume, for instance, that σ_{LR}/σ_{LL} or σ_{RR}/σ_{LL} is bigger than 10^{-6}, then according to the values given by Olive, et al.[6], we are allowed to have only two four-component neutrinos. Since we already know three of them (assuming that all of them are Dirac neutrinos) ν_e, ν_μ and ν_τ, there is not enough room for the third one. Therefore the ratio must be smaller than 10^{-6}.

Using this restriction we find the upper limits for ζ and δ. They are in Table 1, together with the bounds found by Bég et al. using low energy experiments and with the bounds expected in the next round of experiments.[7]

There is another possibility in which we have the tau's neutrino, ν_τ, very massive and hence non-relativistic during the lepton epoch; then its contribution to ρ is negligible. In this case, however, the ν_τ must decay. According to Schramm's calculations, the lifetime of the ν_τ must be 10^{-4} sec. if its mass is around 100 Mev.[8] But if this happens to be true, then the ν_τ has to be a Majorana particle, since the decay necessarily violates lepton number. The problem of Majorana particles present in the early universe are considered in the context of L.R.S. theories, by Roncadelli, et al. (1981).[9]

III. CONCLUSION

There is a greater restriction on the $SU(2)_L \times SU(2)_R \times U(1)$ parameters coming from cosmological arguments than from low energy experiments. The ν_R must decouple early enough in order to avoid any modification to the energy density and consequently to the H_e^4 ratio. Then the ν_R interaction must be at least 10^6 weaker than the ν_L interaction. This constraint imposes severe restrictions on the parameters of the model. We have that $M_{W_R} \geq 30 M_{W_L}$ † and $\tan\zeta < 2\times 10^{-3}$ which are one order of magnitude better bounds than the present experimental limits. The kind of argument used here to constrain the model under consideration could also serve to constrain other models which might require many new particles. If in future experiments a right-handed gauge boson is found with mass lower than 2.5 Tev, this could be an indication that the neutrinos are Majorana eigenstates.

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† Beall, et al. [10] predict $M_R > 1.6$ Tev from the K_L-K_S mass difference calculated using the L-R-S model.

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FIGURE 1

Feynman diagrams which contribute to heat up the right-handed neutrinos.

Table 1. Bounds on the parameters of the model coming from different sources.

	δ	$\tan\zeta$
Bég, et al.	.13	$^{+.054}_{-.06}$
Expt. prediction	.02	±.02
Cosmology	.001	±.002

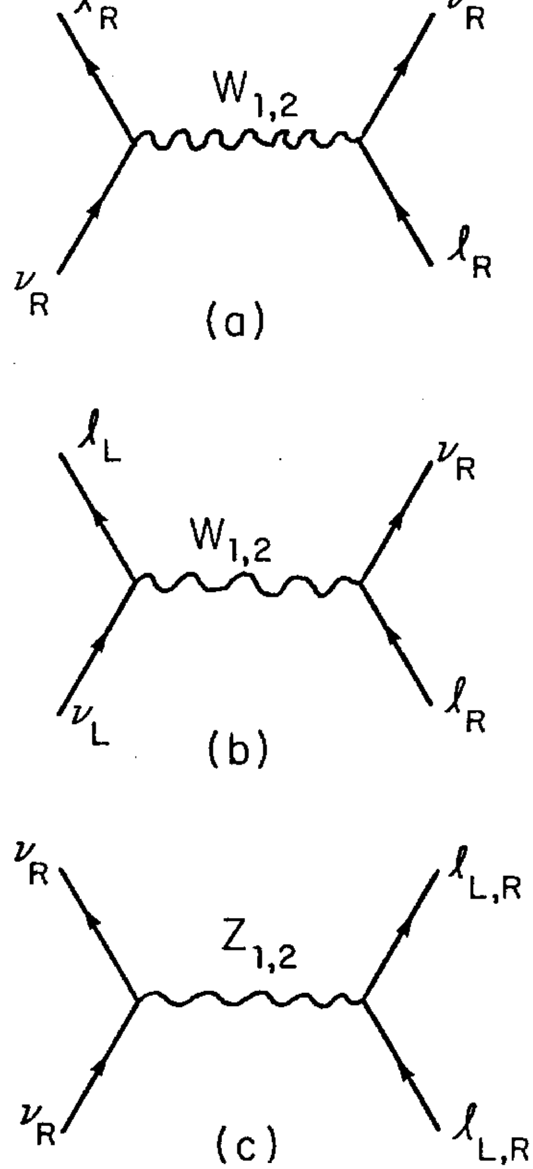
ℓ_R
ν_R
$W_{1,2}$
ν_R
ℓ_R
(a)
ℓ_L
ν_R
$W_{1,2}$
ν_L
ℓ_R
(b)
ν_R
$\ell_{L,R}$
$Z_{1,2}$
ν_R
$\ell_{L,R}$
(c)

Fig. 1