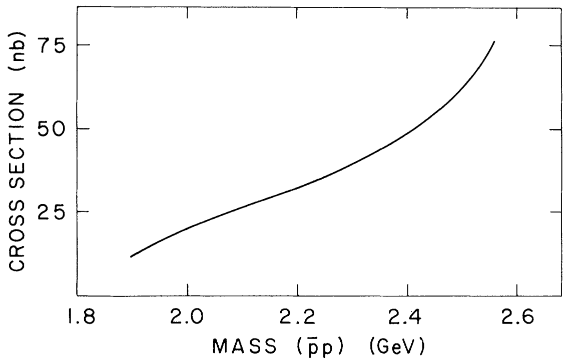

FIG. 4. 95% confidence limits for the production cross section of narrow mass $\bar{p}p$ states.

enhanced production expected in π^+p reactions. Figure 4 indicates the 95%-confidence-level cross-section limit over our mass range.[11]

In summary, we find no evidence for production of narrow $\bar{p}p$ states produced in $\pi^+p \rightarrow \Delta_f^{++}\bar{p}p$. In particular, we are unable to confirm the existence of states at 2.02 and 2.20 GeV/c^2 reported in $\pi^-p \rightarrow \Delta_f^0\bar{p}p$. Those states would appear as >5-standard-deviation effects in our experiment.

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Neutrino Mass and Spontaneous Parity Nonconservation

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In weak-interaction models with spontaneous parity nonconservation, based on the gauge group $SU(2)_L \otimes SU(2)_R \otimes U(1)$, we obtain the following formula for the neutrino mass: $m_{\nu_e} \simeq m_e^2/gm_{W_R}$, where W_R is the gauge boson which mediates right-handed weak interactions. This formula, valid for each lepton generation, relates the maximality of observed parity nonconservation at low energies to the smallness of neutrino masses.

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It is attractive to suppose that observed parity nonconservation in weak interactions is only a low-energy phenomenon, which ought to disappear at high energies. This idea has been implemented in unified gauge theories of electroweak interactions based on the gauge group $SU(2)_L \otimes SU(2)_R$

$\otimes$ U(1),[1] where parity nonconservation arises from the spontaneous breaking of the gauge symmetry. The suppression of the right-handed weak currents in this model owes its origin to the large mass of the right-handed gauge bosons. As far as the structure of the neutrino neutral-current interactions[2] and the parity-nonconserving electron-hadron weak interactions[3] are concerned, this model is indistinguishable from the standard $SU(2)_L \otimes U(1)$ model[4] at the present level of experimental accuracy. There exists, however, a fundamental distinction between the left-right-symmetric models and the pure left-handed $SU(2)_L \otimes U(1)$ models: In the former the neutrino has an arbitrary but finite mass, whereas in the latter it is massless. It is therefore important to understand the smallness of the neutrino mass in left-right symmetric models. Furthermore, it is very suggestive in the context of these models that there may be a connection between the smallness of the neutrino mass and the suppression of the right-handed weak interactions. In this Letter, we propose a model of spontaneous parity nonconservation based on the $SU(2)_L \otimes SU(2)_R \otimes U(1)$ gauge group, where this connection is brought out explicitly. We obtain the following estimate that relates the neutrino mass to the mass of the right-handed gauge bosons (see below for the detailed nature of the approximations):

$$m_{\nu_e} \simeq m_e^2/g m_{W_R}. \tag{1}$$

A similar formula holds for leptons in each generation. This formula is very illuminating in the sense that, in the limit of $m_{W_R} \to \infty$, the neutrino mass goes to zero and we have at the same time a pure $V-A$ theory of weak interactions.

We now proceed to derive Eq. (1) for one generation of leptons and repeat the same procedure for each generation. The main new ingredient of our proposal is that we start with two Majorana[5,6] neutrinos ν and N and choose the left- and right-handed lepton multiplets prior to spontaneous breakdown to be

$$\psi_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix}, \quad \psi_R = \begin{pmatrix} N_R \\ e_R \end{pmatrix}, \tag{2}$$

with $SU(2)_L \otimes SU(2)_R \otimes U(1)$ representation numbers $(\frac{1}{2}, 0, -1)$ and $(0, \frac{1}{2}, -1)$, respectively. The quarks are assigned to left-right doublets as before.[1] We impose the left-right symmetry on the Lagrangian; under this symmetry $\psi_L \leftrightarrow \psi_R$ and this demands that at the tree level, $g_L = g_R$. We now introduce the Higgs multiplets[7] to break the gauge symmetry down to $U(1)_{\rm em}$: φ transforms as the $(\frac{1}{2}, \frac{1}{2}, 0)$ representation of the gauge group; $\Delta_L \equiv (1, 0, 2)$ and $\Delta_R \equiv (0, 1, 2)$. Under left-right discrete symmetry $\varphi \leftrightarrow \varphi^\dagger$ and $\Delta_L \leftrightarrow \Delta_R$. We have already shown[8] earlier that, starting with a left-right-symmetric potential, it is possible to find a domain of coupling parameters in the theory for which we have

$$\langle \varphi \rangle = \begin{pmatrix} \kappa & 0 \\ 0 & \kappa' \end{pmatrix}, \quad \langle \Delta_L \rangle = 0, \quad \langle \Delta_R \rangle = \begin{pmatrix} 0 \\ 0 \\ v \end{pmatrix}. \tag{3}$$

It is easy to see that for $v \gg \kappa', \kappa$, after the first stage of the symmetry breakdown, the local symmetry group is reduced to $SU(2)_L \otimes U(1)$, where U(1) corresponds to $T_{3R} + Y$, which is finally broken down to $U(1)_{\rm em}$ by $\langle \varphi \rangle \neq 0$.

We now proceed to discuss the main result of our paper, i.e., calculation of the neutrino masses. For simplicity, we also assume that $\kappa' \ll \kappa$. The gauge-invariant Yukawa couplings can be written as

$$\mathcal{L}_Y = h_1 \bar{\psi}_L \varphi \psi_R + h_2 \psi_L \tilde{\varphi} \psi_R + h_3 (\psi_L{}^T C i \tau_2 \Delta_L \psi_L + \psi_R{}^T C i \tau_2 \Delta_R \psi_R) + \text{H.c.}, \tag{4}$$

where $\tilde{\varphi} \equiv \tau_2 \varphi^* \tau_2$ and

$$\Delta_{L,R} = \begin{pmatrix} \frac{1}{\sqrt{2}}\delta^+ & \delta^{++} \\ \delta^0 & -\frac{1}{\sqrt{2}}\delta^+ \end{pmatrix}_{L,R},$$

and C is the Dirac charge-conjugation matrix. From (3) and (4) we find for the electron mass

$$m_e \simeq h_2 \kappa \tag{5}$$

and the mass matrix[9] for the ν-N sector is

$$\begin{array}{c} \\ \nu \\ N \end{array} \begin{array}{c} \begin{array}{cc} \bar{\nu} & \bar{N} \end{array} \\ \begin{pmatrix} 0 & h_1\kappa \\ h_1\kappa & h_3 v \end{pmatrix} \end{array}, \tag{6}$$

where we have used the property of Majorana particles $\nu^c = \nu$, $N^c = N$ in showing that $N_R{}^T C N_R$ and $\nu_L{}^T C \nu_L$ are mass terms. We further assume that Yukawa couplings h_1 and h_2 are of the same order of magnitude, i.e., $h_1 \simeq h_2$. It then follows from (6) that $m_N \simeq h_3 v$ and

$$m_\nu = (h_1\kappa)^2/m_N = g m_e^2/h_3 m_{W_R}. \tag{7}$$

This is the main result of our paper. Choosing a reasonable value for h_3, e.g., $h_3 \simeq g^2$, we obtain (1). For the second and third generations of leptons, the corresponding formulas are

$$m_{\nu_\mu} \simeq m_\mu^2/g m_{W_R} \text{ and } m_{\nu_\tau} = m_\tau^2/g m_{W_R}, \tag{8}$$

where we assume the heavy Majorana mass to be generation independent and we ignore generation mixings. This admittedly arbitrary assumption is taken only for illustrative purposes; therefore the values for m_{ν_μ} and m_{ν_τ} should be taken with this in mind.

To point out that these formulas lead to quite reasonable upper limits on the neutrino masses, we note that the analysis of charged-[10] and neutral-current[2] phenomena puts a lower bound on the right-handed W-boson mass, i.e., $m_{W_R} \gtrsim 3m_{W_L}$ (or $m_{W_R} \gtrsim 250-300$ GeV). If we choose $m_{W_R} \gtrsim 300$ GeV, Eqs. (1) and (8) yield $m_{\nu_e} \lesssim 1.5$ eV, $m_{\nu_\mu} < 56$ keV, and $m_{\nu_\tau} \lesssim 18$ MeV, in accord with the present laboratory experiments.[11]

We now comment on the quark sector of the model. As noted earlier, if we restrict ourselves to one generation, the left and right doublets are $(u_L, d_L) \equiv (\frac{1}{2}, 0, \frac{1}{3})$ and $(u_R, d_R) \equiv (0, \frac{1}{2}, \frac{1}{3})$. They obtain their mass through coupling to the Higgs multiplet φ. However, because the U(1) quantum number for quarks is $\frac{1}{3}$, $\Delta_{L,R}$ do not couple to quarks.

We now wish to remark on the following aspects of the model:

(a) At low energies, our model is indistinguishable from the standard model, since $\langle \Delta_R \rangle \neq 0$ still keeps $Y' \equiv T_{3R} + Y$ and $\vec{T}_L$ unbroken; that is, $\mathrm{SU}(2)_L \otimes \mathrm{U}(1)$ is unbroken after the first stage of symmetry breaking. We obtain the following mass relations between the neutral and charged gauge mesons (in the limit $v^2 \gg \kappa^2 + \kappa'^2$):

$$m_{W_L}^2 \simeq m_{Z_L}^2 \cos^2\theta_W, \qquad m_{W_R}^2 \simeq m_{Z_R}^2 (2\cos 2\theta_W / \cos^2\theta_W). \tag{9}$$

Note the factor 2 in (9), which is characteristic of the triplet Higgs boson Δ_R which breaks $\mathrm{SU}(2)_R$.

(b) In order to suppress lepton-number–changing processes such as $\mu \to e\gamma$ and $\mu \to 3e$, it turns out to be desirable to make the mixings of electron generation with μ and τ generation as small as possible. We make it zero naturally by imposing the following discrete symmetry on the Lagrangian

$$\begin{aligned} &\psi_{1L} \to -\psi_{1L}, \quad \psi_{1R} \to -\psi_{1R}, \\ &\psi_{2L} \to \psi_{2L}, \quad \psi_{2R} \to \psi_{2R}, \\ &\psi_{3L} \to \psi_{3L}, \quad \psi_{3R} \to \psi_{3R}, \\ &\varphi \to \varphi, \quad \Delta_L \to \Delta_L, \quad \Delta_R \to \Delta_R, \end{aligned} \tag{10}$$

where $i = 1,2,3$ counts the electron, the muon, and the τ generation, respectively, i.e.,

$$\psi_{1L} = \begin{pmatrix} \nu_e \\ e \end{pmatrix}_L, \quad \psi_{2L} = \begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix}_L,$$

etc. This symmetry obviously forbids e-μ and e-τ mixings and, since it is unbroken, that will be true to all orders in perturbation theory. Therefore, in this model $\mu \to e\gamma$ and $\mu \to 3e$ are forbidden processes. By the same token, neutrino oscillations $\nu_e \leftrightarrow \nu_\mu$, $\nu_e \leftrightarrow \nu_\tau$ are absent in this model (but not $\nu_\mu \leftrightarrow \nu_\tau$). The interesting physical consequences when the above symmetry is relaxed will be dealt with elsewhere.

(c) A further characteristic of our model is the existence of doubly charged Higgs bosons δ_L^{++} and δ_R^{++}. However, it is easily seen that their masses are of order m_{W_R} and therefore they are not expected to play an important role at low energies.

(d) The presence of Majorana neutrinos will allow for neutrinoless double β decay.[12,13] The contribution coming from the exchanges of W_L involves the light Majorana neutrino ν_e and is known[13] to be a few orders of magnitude below the experimentally allowed value. However, it is possible to exchange W_R's in which case the process goes through the heavy Majorana particle N as an internal state. We just mention that in this case one obtains the limit on N using the analysis of Ref. 13: $m_N \gtrsim (m_{W_L}/m_{W_R})^4 \times 10^4$ GeV $\gtrsim 10^2$ GeV, which is definitely satisfied, since $m_N \simeq m_{W_R} \simeq 300$ GeV. It should be emphasized, though, that more precise measurements of double β decay could in principle provide a more stringent lower bound on m_N, or in turn on m_{W_R}.

In summary, we have constructed a realistic and simple model with spontaneous parity nonconservation, where the suppression of $V+A$ currents is proportional to neutrino mass. The model provides, therefore, an understanding[14] of a tiny neutrino mass. We believe that it makes the search for the effects due to finite m_{W_R} even more warranted than before.

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