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Neutrino masses, neutral vector bosons and the scale of $B-L$ breaking

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Extensions of the standard model with extra neutral vector bosons predict additional fermions whose allowed masses are constrained by the masses of the new vector bosons. For the minimal case with one extra Z' boson and an extended gauge group contained in the unified group SO(10), we derive upper and lower bounds on the masses of heavy Majorana neutrinos. We also discuss implications for the standard model Higgs boson.

Despite its extraordinary success the standard model of strong and electroweak interactions is generally believed to be only the low energy approximation of a more fundamental, unified theory. A possible remnant of unification at energies much below the unification scale are extra neutral vector bosons, which in recent years have been discussed in particular in connection with superstring theories [1]. Extended gauge theories, which contain U(1) factors in addition to the standard model gauge group, also predict fermions in addition to the quarks and leptons of the standard model. They are contained in the possible anomaly free fermion representations of the unified theory. If one or more U(1) factors of the unified group are spontaneously broken far below the unification scale, also some of these fermions will in general acquire masses much smaller than the unification mass.

In the following we shall consider the minimal extension of the standard model of this type, which is based on the gauge group $SU(3)_C\times SU(2)_L\times U(1)_Y\times U(1)_{Y'}$ contained in the unified group SO(10). The smallest complex SO(10) representation, the 16 plet, contains a "right-handed neutrino" ν_R in addition to the 15 Weyl fermions of one quark-lepton family of the standard model. ν_R is a singlet with respect to the standard model gauge group. Hence, in the sequential breaking

$$SU(3)_C\times SU(2)_L\times U(1)_Y\times U(1)_{Y'} \xrightarrow{v'} SU(3)_C\times SU(2)_L\times U(1)_Y \xrightarrow{v} SU(3)_C\times U(1)_{em}, \tag{1}$$

ν_R can acquire a Majorana mass of order v' and thereby explain the smallness of the ν_e, ν_μ and ν_τ masses via the see-saw mechanism [2]. Here v' and v are the vacuum expectation values which break the subgroups $(U)_{Y'}$ and $SU(2)_L\times U(1)_Y$, and $U(1)_{em}$ is the unbroken group of electromagnetic interactions. As we shall see, the Majorana masses of the heavy neutrinos are bounded from above essentially by v', and they are also bounded from below by the experimental upper bounds on the light neutrino masses.

Heavy Majorana neutrinos with masses up to a few hundred GeV can be produced at present and future ep and e^+e^- colliders, and detailed discovery limits have recently been worked out [3]. In order to obtain the precise relation between the Z′ vector boson mass and the masses of the heavy Majorana neutrinos we study in the following the spontaneous breaking of the extended gauge symmetry. Contrary to the fermion content, which is constrained by the requirement of anomaly freedom, there are no restrictions on the Higgs fields which remain light compared to the unification scale. We shall therefore begin with the minimal Higgs sector necessary to break the extended gauge symmetry and comment later on the effects of additional Higgs fields.

The spontaneous breaking of $U(1)_{Y'}$ requires a complex scalar field χ in addition to the doublet ϕ of Higgs fields, and the lagrangian, which also yields Majorana masses for the right-handed neutrinos, reads

$$\begin{aligned}\mathscr{L}= &-\tfrac{1}{4}W^I_{\mu\nu}W^{I\mu\nu}-\tfrac{1}{4}B_{\mu\nu}B^{\mu\nu}-\tfrac{1}{4}C_{\mu\nu}C^{\mu\nu}\\ &+\mathrm{i}\bar{\ell}\not{D}\ell+\mathrm{i}\overline{\nu_R}\not{D}\nu_R+\mathrm{i}\overline{e_R}\not{D}e_R\\ &+(D_\mu\phi)^\dagger(D^\mu\phi)+(D_\mu\chi)^\dagger(D^\mu\chi)-V(\chi,\phi)\\ &-\bar{\ell}\phi g_\nu\nu_R-\overline{\nu_R}\,g_\nu^\dagger\phi^\dagger\ell-\bar{\ell}\tilde{\phi}g_e e_R-\overline{e_R}g_e^\dagger\tilde{\phi}^\dagger\ell\\ &-\tfrac{1}{2}\chi\overline{\nu_R^c}h\nu_R-\tfrac{1}{2}\chi^\dagger\overline{\nu_R}h^\dagger\nu_R^c\,. \end{aligned}\tag{2}$$

$\phi=(\varphi^0,\varphi^-)$, $\tilde{\phi}=(\varphi^-,-\varphi^0)$ and χ are the SU(2) doublet and singlet Higgs fields, $\ell=(\nu_L, e_L^-)$ and ν_R are the lepton doublet and the right-handed neutrino. g_ν and h are 3×3 complex Yukawa matrices with $h=h^T$. The covariant derivatives differ from the standard model expressions by the term $-\mathrm{i}\sqrt{\tfrac{3}{5}}\,g'\times Y'C\mu$, with $Y'(\ell)=\tfrac{3}{4}$, $Y'(e_R)=\tfrac{1}{4}$, $Y'(\nu_R)=\tfrac{5}{4}$, $Y'(\phi)=-\tfrac{1}{2}$, $Y'(\chi)=-\tfrac{5}{2}$. All terms involving quark fields have been omitted. B_μ, C_μ and W^I_μ are the $U(1)_Y$, $U(1)_{Y'}$ and $SU(2)_L$ gauge fields, and $B_{\mu\nu}$, $C_{\mu\nu}$ and $W^I_{\mu\nu}$ are the corresponding field strengths. The couplings of the vector field C_μ are determined by the orthonormality conditions for the U(1) charges $[Y'=aY+b(B-L)]$

$$\mathrm{tr}[YY']=0\,,\quad \mathrm{tr}[Y^2]=\tfrac{3}{5}\,\mathrm{tr}[Y'^2]\,,\tag{3}$$

where the trace extends over all fermions, which yields

$$Y'=Y-\tfrac{5}{4}(B-L)\,.\tag{4}$$

The gauge coupling constants of B_μ and C_μ are equal as the β-functions of both U(1) groups are the same. The most general Higgs potential is given by

$$\begin{aligned}V(\chi,\phi)=&\mu_1\phi^\dagger\phi+\mu_2\chi^\dagger\chi+\tfrac{1}{2}\lambda_1(\phi^\dagger\phi)^2\\ &+\tfrac{1}{2}\lambda_2(\chi^\dagger\chi)^2+\lambda_3(\chi^\dagger\chi)(\phi^\dagger\phi)\,,\end{aligned}\tag{5}$$

where $\lambda_3>-(\lambda_1\lambda_2)^{1/2}$, so that the potential is bounded from below. This Higgs sector is precisely the well known "singlet majoron" model [4] for spontaneous breaking of lepton number.

It is now straightforward to minimize the Higgs potential (5), to determine the vacuum expectation values v and v' of the Higgs fields ϕ and χ, and to diagonalize the vector boson mass matrix. In addition to the photon one obtains two massive neutral vector bosons Z and Z′,

$$\begin{aligned}Z_\mu&=\cos\xi\,Z^0_\mu-\sin\xi\,C_\mu\,,\\ Z'_\mu&=\sin\xi\,Z^0_\mu+\cos\xi\,C_\mu\,,\end{aligned}\tag{6}$$

where $Z^0_\mu=-\sin\Theta\,B_\mu+\cos\Theta\,W^3_\mu$ corresponds to the neutral vector boson of the standard model. Neglecting fermion masses, predictions of the standard model can be expressed in terms of three independent parameters which one may choose as the fine structure constant α, the Fermi constant G_F and the Z boson mass m_Z. Since in our model the Higgs sector is specified, there is only one additional parameter which we choose to be the Z′ boson mass $m_{Z'}$. The mixing angles Θ and ξ can now be expressed in terms of α, G_F, m_Z and $m_{Z'}$. One then obtains the following exact relations ($\mu^2=\pi\alpha/\sqrt{2}\,G_F$):

$$\sin^2\Theta=\frac{1}{2}-\left(\frac{1}{4}-\frac{\mu^2}{\rho m_Z^2}\right)^{1/2},\tag{7}$$

$$\sin 2\xi=-2\sqrt{\tfrac{3}{5}}\sin\Theta\frac{\rho m_Z^2}{m_{Z'}^2-m_Z^2}\,,\tag{8}$$

$$\Delta\rho\equiv\rho-1=\sin^2\xi\frac{m_{Z'}^2-m_Z^2}{m_Z^2}\,.\tag{9}$$

Here we have introduced as auxiliary quantity the parameter ρ, which is related to the ratio of neutral to charged current amplitudes at zero momentum transfer. In the relevant case $m_Z^2\ll m_{Z'}^2$ the quantities $\sin\Theta$, $\sin 2\xi$ and $\Delta\rho$ may be expanded around the standard model limit $m_{Z'}=\infty$, and one obtains to first order in $m_Z^2/m_{Z'}^2$:

$$\xi=-\sqrt{\tfrac{3}{5}}\sin\hat{\Theta}\frac{m_Z^2}{m_{Z'}^2}\,,\tag{10}$$

$$\Delta\rho = \tfrac{2}{3}\sin^2\hat{\Theta}\frac{m_Z^2}{m_{Z'}^2}, \tag{11}$$

$$\sin\Theta = \sin\hat{\Theta}\left(1 - \frac{1}{12}\frac{\sin^2\hat{\Theta}}{\cos 2\hat{\Theta}}\frac{m_Z^2}{m_{Z'}^2}\right), \tag{12}$$

where $\hat{\Theta} = \Theta|_{\rho=1}$.

The mixing angles Θ and ξ determine the couplings of the vector bosons Z and Z′ to quarks and leptons. The interaction lagrangian is given by

$$\mathscr{L} = J^{\mu}_{NC} Z_\mu + J'^{\mu}_{NC} Z'_\mu, \tag{13}$$

where

$$J^{\mu}_{NC} = \frac{g\cos\xi}{2\cos\Theta}\sum_i \bar{\psi}_i\gamma^\mu(V^i - A^i\gamma_5)\psi_i, \tag{14}$$

$$V^i = T_{3L}(i) - 2\sin^2\Theta\, Q_i$$
$$+\sqrt{\tfrac{2}{3}}\tan\xi\sin\Theta\,[T_{3L}(i) - 2Q_i + \tfrac{5}{2}(B-L)_i], \tag{15}$$

$$A^i = (1 + \sqrt{\tfrac{2}{3}}\tan\xi\sin\Theta)T_{3L}(i), \tag{16}$$

$$J'^{\mu}_{NC} = \frac{g\sin\xi}{2\cos\Theta}\sum_i \bar{\psi}_i\gamma^\mu(V'^i - A'^i\gamma_5)\psi_i, \tag{17}$$

$$V'^i = T_{3L}(i) - 2\sin^2\Theta\, Q_i$$
$$-\sqrt{\tfrac{2}{3}}\,\mathrm{ctg}\,\xi\sin\Theta\,[T_{3L}(i) - 2Q_i + \tfrac{5}{2}(B-L)_i], \tag{18}$$

$$A'^i = (1 - \sqrt{\tfrac{2}{3}}\,\mathrm{ctg}\,\xi\sin\Theta)T_{3L}(i). \tag{19}$$

Here the sum is over all quarks and leptons, and $T_{3L}(i)$, Q_i and $(B-L)_i$ denote weak isospin, electric charge and $B-L$ of the ith fermion.

The Z–Z′ mixing affects various observables which have been measured with high precision at LEP [5]. Two particularly sensitive quantities are the ρ-parameter [cf. eqs. (9) and (11)] and the leptonic width of the Z boson, given by

$$\Gamma_{ee} = \frac{G_F m_Z^3\rho}{6\pi\sqrt{2}}\,[\tfrac{1}{2} - 2\sin^2\Theta + 4\sin^4\Theta$$
$$+\xi\sqrt{\tfrac{2}{3}}\sin\Theta(\tfrac{3}{2} - 4\sin^2\Theta)]. \tag{20}$$

The three variables Θ, ρ and ξ on the right hand side of eq. (20) can be expressed in terms of the independent parameters α, G_F, m_Z and $m_{Z'}$ according to eqs. (10)–(12). Here we neglect the effect of the mixing between heavy and light neutrinos on the Fermi constant G_F. A detailed comparison of LEP data with the standard model predictions shows that the mixing angle $|\xi|$ has to be smaller than about 10^{-2} (cf. refs. [6–8]). From eqs. (11) and (20) one finds that this corresponds to an accuracy of about 0.4% in ρ and Γ_{ee}. Eq. (10) then yields the lower bound on the Z′ boson mass ($|\xi| < 10^{-2}$):

$$m_{Z'} = \sqrt{\sqrt{\tfrac{2}{3}}\frac{\sin\hat{\Theta}}{|\xi|}}\, m_Z > 570\ \text{GeV}, \tag{21}$$

where we have used $\sin^2\hat{\Theta} = 0.23$ and $m_Z = 91$ GeV.

In eqs. (14)–(19) the neutral currents J^{μ}_{NC} and J'^{μ}_{NC} are expressed in terms of weak eigenstates which are related to mass eigenstates by unitary transformations. For the neutrinos these transformations are determined by requiring that the two Majorana mass matrices

$$m = hv', \tag{22}$$

$$m_\nu = -m_D\frac{1}{m}m_D^T, \tag{23}$$

with

$$m_D = g_\nu v, \tag{24}$$

are diagonal and real (cf., e.g., ref. [3]). g_ν is the 3×3 matrix of neutrino Yukawa couplings. The eigenvalues of m are the masses of three heavy Majorana neutrinos, the eigenvalues of m_ν correspond to the masses of ν_e, ν_μ and ν_τ.

Using eqs. (22), (23) and the relation [cf. eq. (2)]

$$m_{Z'}^2 = \tfrac{25}{3}g'^2v'^2[1 + O(v^2/v'^2)], \tag{25}$$

we can now estimate the mass range, in which heavy Majorana neutrinos are to be expected. We assume that the Yukawa couplings g_ν and h vary between the smallest and largest known Yukawa couplings $g_e = m_e/v = 3\times 10^{-6}$ and $g_t \approx 1$, which is about the largest value allowed by "triviality" considerations [9]. From eqs. (22) and (25) we then obtain an upper bound on the Majorana masses m_i. Furthermore eq. (23) and the experimental upper limits [10] $m_{\nu_e} < 17$ eV, $m_{\nu_\mu} < 0.27$ MeV, $m_{\nu_\tau} < 35$ MeV imply the lower bound $\det(m) \equiv m_1 m_2 m_3 > m_e^6/m_{\nu_e}m_{\nu_\mu}m_{\nu_\tau}$. Upper and lower bound together yield the allowed mass range

$$5\ \text{MeV} < (m_1 m_2 m_3)^{1/3}, \tag{26}$$

$$m_i < 970\,\mathrm{GeV}\left(\frac{m_{Z'}}{1\,\mathrm{TeV}}\right), \quad i=1,2,3\,. \tag{27}$$

If, instead of m_e^3, we take $\det(m_\ell) \equiv m_e m_\mu m_\tau$ as lower bound on the determinant of the Dirac neutrino mass matrix m_D, then the lower bound on $(m_1 m_2 m_3)^{1/3}$ increases to ~ 40 GeV. The lower bound crucially depends on the assumption about the Yukawa couplings g_ν; it disappears for $g_\nu = 0$.

From eqs. (26) and (27) we conclude that Majorana neutrinos with masses of a few tens of GeV or a few hundred GeV are likely to exist if an SO(10) Z′ boson with mass of order 1 TeV is found. If the mixings of these heavy neutrinos with ν_e, ν_μ and ν_τ are not too small, they can be produced in charged and neutral current processes at HERA, LEP or LEP⊗LHC [3].

The upper bound on the heavy neutrino masses stems from the fact that the vacuum expectation value, which generates the Majorana mass term, unavoidably also contributes to the Z′ vector boson mass. Hence, the upper bound (27) on m_i, which we obtained for the minimal Higgs sector, will become even more stringent if more Higgs scalars are introduced which contribute to the Z′ boson mass, but do not couple to $\overline{\nu_R}\nu_R^c$. In the case of more than one singlet scalar field with the quantum numbers of χ our assumption on the maximal Yukawa coupling refers to the linear combination of scalar fields which acquires a vacuum expectation value.

One may worry that a low mass scale of $B-L$ breaking of order 1 TeV is in conflict with constraints on lepton number violating processes from low energy experiments [11]. This is not the case although some bounds on mixing angles in charged and neutral currents can be derived. As an example, let us briefly mention neutrinoless double β decay, which provides the most stringent bound on lepton number violation. The present experimental bound on light Majorana neutrino masses reads [12]

$$\langle m_{\nu_e}\rangle = \sum_{\ell=e,\mu,\tau} U_{e\nu_\ell}^2 m_{\nu_\ell} < (1\text{–}10)\ \mathrm{eV}\,, \tag{28}$$

where $U_{e\nu_\ell}$ denotes elements of the leptonic Kobayashi–Maskawa (MK) matrix. For instance, with $m_{\nu_e} = m_e^2/m$ and $m \sim 500$ GeV one has $m_{\nu_e} \sim 0.5$ eV, which satisfies the bound (28). Furthermore, if $U_{e\nu_\ell} = O(\sqrt{m_{\nu_e}/m_{\nu_\ell}})$, which approximately holds for quark KM matrix elements, also the contributions from ν_μ and ν_τ satisfy the bound (28). A mass scale of $B-L$ breaking of order 1 TeV is also cosmologically acceptable [13,14].

Let us finally consider the Higgs sector. The fields ϕ and χ contain the neutral bosons H^0 and H'^0:

$$\phi^\dagger\phi = \left(v + \frac{1}{\sqrt{2}}H^0\right)^2,$$

$$\chi^\dagger\chi = \left(v' + \frac{1}{\sqrt{2}}H'^0\right)^2. \tag{29}$$

The mass eigenstates H and H' are linear combinations of H^0 and H'^0, i.e., $H^0 = \cos\beta H + \sin\beta H'$ and $H'^0 = -\sin\beta H + \cos\beta H'$. For the mixing angle one finds

$$\tan 2\beta = \frac{2\lambda_3 v v'}{\lambda_2 v'^2 - \lambda_1 v^2}\,. \tag{30}$$

The corresponding Higgs boson masses are

$$m_H^2 = \lambda_1 v^2 + \lambda_2 v'^2 - \sqrt{(\lambda_1 v^2 - \lambda_2 v'^2)^2 + 4\lambda_3^2 v^2 v'^2}\,, \tag{31}$$

$$m_{H'}^2 = \lambda_1 v^2 + \lambda_2 v'^2 + \sqrt{(\lambda_1 v^2 - \lambda_2 v'^2)^2 + 4\lambda_3^2 v^2 v'^2}\,. \tag{32}$$

For $v' = O(1\ \mathrm{TeV})$ the mixing angle β and the splitting between the two Higgs boson masses may be large. Furthermore the couplings of one Higgs boson to ordinary quarks and leptons could be substantially reduced. Since H'^0 interacts only with C_μ, ν_R and H^0, the couplings of the mass eigenstates H and H' to W and Z bosons, quarks and charged leptons have essentially the same structure as the corresponding Higgs couplings in the standard model, however, their size is reduced by $\cos\beta$ and $\sin\beta$, respectively. In the extreme case $\beta = \frac{1}{2}\pi$ the couplings of the lighter Higgs boson to quarks and leptons would even vanish!

To summarize, we have shown that the existence of a neutral Z′ vector boson with mass of order 1 TeV has important implications for other particles in the low energy theory. For an extended gauge group contained in SO(10) we have estimated upper and lower bounds on the masses of heavy Majorana neutrinos, and we have studied implications for the standard model Higgs boson. For Z′ vector bosons contained in E_6 similar bounds on masses of the additional neutral and charged leptons and quarks can be derived.

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