Decays of weak vector bosons and t quarks into doubly charged Higgs scalars

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We evaluate the decays of W^+ and Z^0 bosons into χ^{++}, the doubly charged member of the scalar triplet in the Gelmini-Roncadelli model in which spontaneously broken $B-L$ symmetry gives Majorana masses to neutrinos. For $M(\chi^{++})<28$ GeV, the branching fractions exceed 4% for $W^+\to\chi^{++}\chi^-$ and 2% for $Z^0\to\chi^{++}\chi^{--}$. Sequential $\chi^{++}\to l_1^+ l_2^+$ leptonic decays could be substantially higher than previously recognized and could be dominant. Decay distributions for leptons of this origin are calculated for $\bar{p}p$ colliders. We also evaluate $t\to b\chi^{++}\chi^-$ decay.

A possible origin of Majorana neutrino mass is the spontaneous breakdown of $B-L$ global symmetry in the standard SU(2)×U(1) gauge model, as proposed recently by Gelmini and Roncadelli[1] (GR). This breakdown occurs via a Higgs scalar triplet whose neutral member acquires a vacuum expectation value v_T much smaller than that of the usual Higgs doublet v_D. In this model the physical scalars and their masses are as follows[1,2]: χ^{++} of mass M (> 15 GeV from e^+e^- experiments), χ^+ of mass $M/\sqrt{2}$, a light χ^0 of mass $m\sim v_T<100$ keV, a massless Nambu-Goldstone boson M^0, and the usual Higgs scalar H^0 of the standard model. The $W^\pm$ and Z^0 have gauge couplings to bilinear combinations of the new scalars, which may therefore appear as decay products. The new heavy scalars themselves decay either by gauge couplings or by their Yukawa couplings to leptons. For doubly charged scalars, these lepton couplings provide spectacular signatures for detection if the branching fractions are large enough. In calculating branching fractions for the intial stages $W^+\to\chi^{++}\chi^-$ or $Z^0\to\chi^{++}\chi^{--}$, the only uncertainty is the value of M. In the sequential decay $\chi^{++}\to l_1^+ l_2^+$, the branching fraction depends critically on the Yukawa coupling strengths to leptons, for which some upper limits are known.[2–5]

In the present work we evaluate branching fractions for these two decay stages as functions of M and the Yukawa coupling to leptons. We find that the $W^+\to\chi^{++}\chi^-$ and $Z^0\to\chi^{++}\chi^{--}$ modes are of order several percent for a substantial allowed range of M; the sequential $\chi^{++}\to l_1^+ l_2^+$ branching fractions are larger than previously recognized, and the readily measurable e^+e^+, $\mu^+\mu^+$, and $e^+\mu^+$ modes could give as much as $\frac{4}{9}$ of χ^{++} decays within present coupling limits. Finally, we calculate some useful distributions of leptons from the $W(Z)\to\chi^{++}\to l_1^+ l_2^+$ decay chain, for W and Z produced at a $\bar{p}p$ collider with $\sqrt{s}=2$ TeV. We also evaluate the widths for $t\to b\chi^{++}\chi^-$ decay.

The gauge couplings of W and Z to the new scalars of the GR model[1] are given by the Lagrangian

$$\mathscr{L}=-\frac{ig}{\cos\theta_W}[(1-2x_W)\chi^{--}\overleftrightarrow{\partial}_\mu\chi^{++}+x_W\chi^{-}\overleftrightarrow{\partial}_\mu\chi^{+}-i\chi^0\overleftrightarrow{\partial}_\mu M^0]Z_\mu -ig[\chi^{+}\overleftrightarrow{\partial}_\mu\chi^{--}+(2)^{-1/2}(\chi^0-iM^0)\overleftrightarrow{\partial}_\mu\chi^-]W_\mu^+ \,, \quad (1)$$

where $g^2=4\sqrt{2}G_F M_W{}^2$ and $x_W=\sin^2\theta_W$. Beside the standard model decay modes, new scalars give rise to $W^+\to\chi^{++}\chi^-$, $\chi^+\chi^0$, χ^+M^0, and $Z^0\to\chi^{++}\chi^{--}$, $\chi^+\chi^-$, $\chi^0 M$ decays. The branching fractions for these new modes have been given in Ref. 2 in the limit of light M only. We calculate these branching fractions without approximation as functions of M, taking $M_Z=92$ GeV, $M_W=82$ GeV, $\sin^2\theta_W=0.21$ and the usual three generations of fermions (with t-quark mass of 20 GeV). The partial widths into χ states are

$$\begin{aligned}
&\Gamma(W^+\to\chi^{++}\chi^-)=\Gamma_W^0(1-3w^2+w^4/4)^{3/2}\ ,\\
&\Gamma(W^+\to\chi^+M^0)=\Gamma(W^+\to\chi^+\chi^0)\\
&\qquad=\tfrac{1}{2}\Gamma_W^0(1-w^2/2)^3\ ,\\
&\Gamma(Z\to\chi^{++}\chi^{--})=\Gamma_Z^0(1-2x_W)^2(1-4z^2)^{3/2}\ ,\\
&\Gamma(Z\to\chi^+\chi^-)=\Gamma_Z^0 x_W{}^2(1-2z^2)^{3/2}\ ,\\
&\Gamma(Z\to M^0\chi^0)=\Gamma_Z^0\ ,
\end{aligned} \tag{2}$$

where

$$\Gamma_W^0/M_W{}^3=\Gamma_Z^0/M_Z{}^3=(6\pi)^{-1}(G_F/\sqrt{2})$$

and

$$w=M/M_W,\ \ z=M/M_Z\ .$$

The results for branching fractions are shown in Fig. 1; the phase-space suppression as M increases is evident. Nevertheless, the branching fractions to χ^{++} are high enough to be interesting over a considerable range of M, especially for $M<28$ GeV for which $B(W^+\to\chi^{++}\chi^-)>4\%$ and $B(Z\to\chi^{++}\chi^{--})>2\%$.

The Yukawa couplings of the doubly charged Higgs boson to leptons are given by

$$\mathscr{L}=\frac{1}{2\sqrt{2}}\sum_{l,l'}g_{ll'}[\,\bar{l}^c(1+\gamma_5)l'\chi^{++}+\bar{l}'(1-\gamma_5)l^c\chi^{--}]\ , \tag{3}$$

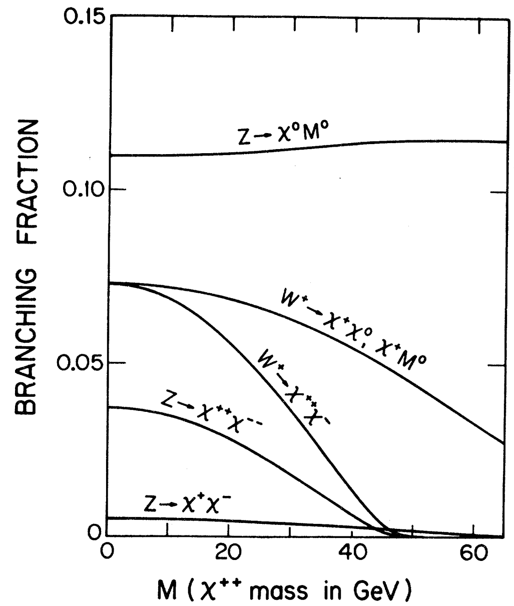

FIG. 1. Predicted branching fractions for W,Z decays to scalars in the GR model is shown versus the χ^{++} mass M.

where superscript c denotes the charge-conjugate lepton field. In Feynman amplitudes a combinatorial factor of 2 usually occurs at the χ^{++} vertex with two leptons.

The χ^{++} decays proceed partly by gauge couplings through a virtual W boson $\chi^{++}\to\chi^+(q\bar{q})^+$, $\chi^+l^+\nu$, and partly by Yukawa couplings to leptons $\chi^{++}\to l^+l'^+$. The branching fractions have been estimated in Ref. 2, but their quoted width for $\chi^{++}\to\chi^+J^+$ omits the very important effect of the χ^+ mass on the phase space integral. The partial width for a typical decay mode of the latter kind is

$$\begin{aligned}
\Gamma(\chi^{++}\to\chi^+l^+\nu)&=(24\pi^3)^{-1}G_F{}^2M^5\int_0^{(3/2)-\sqrt{2}}(\tfrac{1}{4}-3x+x^2)^{3/2}dx\\
&=(24\pi^3)^{-1}G_F{}^2M^5(3/64)(16\ln 2-11)\\
&=5.7\times10^{-6}G_F{}^2M^5
\end{aligned} \tag{4}$$

neglecting lepton masses; the quark channels $\chi^{++}\to\chi^+u\bar{d}$ and $\chi^{++}\to\chi^+c\bar{s}$ each have about three times this width. The variable x in Eq. (4) is the invariant mass squared of the two fermions in units of M^2; the upper limit of x integration exhibits the suppression of phase space by the χ^+ mass. The partial widths for pure leptonic modes are given by

$$\Gamma(\chi^{++}\rightarrow e^+e^+)=g_{ee}{}^2M/(16\pi)\ ,$$
$$\Gamma(\chi^{++}\rightarrow e^+\mu^+)=g_{e\mu}{}^2M/(8\pi)\ , \qquad (5)$$
$$\Gamma(\chi^{+}\rightarrow e^+\nu)=(g^2)_{ee}M/(16\sqrt{2}\pi)$$

with analogous results for other modes, neglecting $m_l{}^2/M^2$. Here g^2 denotes the square of the matrix g, so that $(g^2)_{ee}=\sum_l (g_{el})^2$.

The present upper limits on lepton couplings are

$$g_{ee}\leq 10^{-3}\ \text{(Ref. 2) or}\ \leq 10^{-4}\ \text{(Ref. 5)}\ \beta\beta\ \text{decay}\ , \qquad (6a)$$

$$g_{\mu\mu}<[(g^2)_{\mu\mu}]^{1/2}<1.6\times 10^{-2}\ \text{(Ref. 3)}\ K\rightarrow\mu\ \text{decay}\ , \qquad (6b)$$

$$g_{\mu e}<[(g^2)_{ee}]^{1/2}<7\times 10^{-3}\ \text{(Ref. 3)}\ K\rightarrow l\nu\ \text{universality}\ , \qquad (6c)$$

$$\frac{g_{ee}g_{e\mu}}{M^2}<4G_F[B(\mu\rightarrow ee\bar{e})]^{1/2}<2\times 10^{-9}\ \text{GeV}^{-2}\ , \qquad (6d)$$

and

$$\frac{g_{ee}g_{e\tau}}{M^2}<4G_F[B(\tau\rightarrow ee\bar{e})/B(\tau\rightarrow e\bar{\nu}_e\nu_\tau)]^{1/2}<2\times 10^{-6}\ \text{GeV}^{-2}\ . \qquad (6e)$$

In the latter two entries a difference in the γ matrix traces of $\mu\rightarrow ee\bar{e}$ and $\mu\rightarrow e\nu\bar{\nu}$ and a combinatorial factor in $\mu\rightarrow ee\bar{e}$ are taken into account. The experimental branching-fraction limits

$$B(\mu\rightarrow ee\bar{e})<1.9\times 10^{-9}$$

and

$$B(\tau\rightarrow ee\bar{e})<4\times 10^{-4}$$

have been used.[6] For $g_{e\mu}=g_{ee}=10^{-3}$, Eq. (6d) requires $M>22$ GeV.

If all couplings were approximately equal, two of the three neutrino masses would be nearly zero. With this coupling assumption the readily measurable e^+e^+, $e^+\mu^+$, and $\mu^+\mu^+$ modes together contribute $\frac{4}{9}$ of the leptonic width; the branching fraction for these combined modes is shown versus M in Fig. 2 for the cases $g_{ll'}=10^{-3}$, $g_{ll'}=10^{-4}$, and $g_{ll'}=10^{-5}$. The two-body leptonic modes are dominant for $M<70$ GeV with $g_{ll'}=10^{-3}$.

The quark subprocess cross sections for the production of the χ^{++} in $\bar{p}p$ collisions are given by

$$d\hat{\sigma}(q_i\bar{q}_j\rightarrow W^+\rightarrow \chi^{++}\chi^-)=(24\pi)^{-1}G_F{}^2M_W{}^4\,|\,U_{ij}\,|^2\hat{s}\,|\,P_W\,|^{-2}\lambda^{3/2}(1,M^2/\hat{s},\tfrac{1}{2}M^2/\hat{s})\sin^3\hat{\theta}\,d\hat{\theta}\ ,$$
$$d\hat{\sigma}(q_i\bar{q}_i\rightarrow Z,\gamma\rightarrow \chi^{++}\chi^{--})=(96\pi)^{-1}G_F{}^2M_Z{}^4F\hat{s}\,|\,P_Z\,|^{-2}(1-4M^2/\hat{s})^{3/2}\sin^3\hat{\theta}\,d\hat{\theta}\ , \qquad (7)$$

where $P_\alpha=\hat{s}-M_\alpha{}^2+iM_\alpha\Gamma_\alpha$ is a weak-boson propagator factor, $\Gamma_W\approx\Gamma_Z\approx 3.2$ GeV, U_{ij} is the quark mixing matrix, and

$$F=(1-2x_W)^2+|\,(1-2x_W)(1-4\,|\,e_q\,|\,x_W)+\sqrt{2}\,|\,e_q\,|\,e^2G_F{}^{-1}M_Z{}^{-2}P_Z/\hat{s}\,|^2\ .$$

In these formulas, $\hat{\theta}$ is the c.m. angle of the χ^{++} with respect to the quark q_i and e_q is the quark charge. The cross section for χ^{++} production in $\bar{p}p$ collisions is obtained by folding the above subprocess cross sections with the probability distributions of the constituents in the initial hadrons. Figure 3 shows distributions of μ^+ and $\mu^+\mu^+$ resulting from χ^{++} production and leptonic decay in $\bar{p}p$ collisions at $\sqrt{s}=2$ TeV. The transverse momentum of the $\mu^+\mu^+$ has a prominent peak at 35 GeV which offers a clear signature of this process. The distribution in the azimuthal angle between the two μ^+ in the plane perpendicular to the beam direction has a distinctive peak near 55°.

The t quark has a decay mode $t\rightarrow b\chi^{++}\chi^-$ if the χ^{++} mass M is less than $(2-\sqrt{2})(m_t-m_b)$. The differential decay probability for this mode is

$$d\Gamma/dQ^2=(96\pi^3m_t)^{-1}G_F{}^2\lambda^{1/2}(1,Q^2/m_t{}^2,m_b{}^2/m_t{}^2)\lambda^{1/2}(1,\alpha/Q^2,\tfrac{1}{2}\beta/Q^2)$$
$$\times\{\,[(m_t{}^2-m_b{}^2)^2+(m_t{}^2+m_b{}^2)Q^2-2Q^4]\lambda(1,\alpha/Q^2,\tfrac{1}{2}\beta/Q^2)/(1-Q^2/M_W{}^2)^2$$
$$+3[(m_t{}^2-m_b{}^2)^2-(m_t{}^2+m_b{}^2)Q^2](\alpha-\tfrac{1}{2}\beta)^2/Q^4\,\}\ , \qquad (8)$$

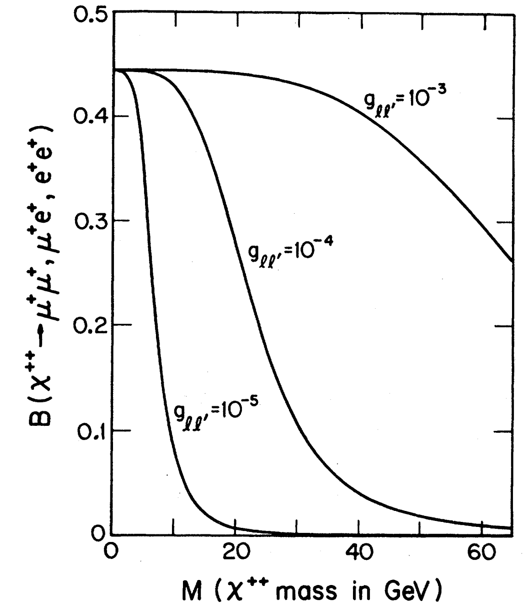

FIG. 2. Combined branching fractions of the X^{++} into e^+e^+, μ^+e^+, $\mu^+\mu^+$ versus the X^{++} mass M, assuming equal couplings $g_{ll'}$. The cases $g_{ll'}=10^{-3}$, 10^{-4}, and 10^{-5} are illustrated.

where Q^2 is the square of the four-momentum transferred to $X^{++}X^-$ by the W boson, and $\alpha=\beta=M^2$. The corresponding decay rate for the summed modes $t \to bX^+M^0$ and $t \to bX^+X^0$ can be obtained with $\alpha=\frac{1}{2}M^2$ and $\beta=0$ in Eq. (8). Figure 4 shows the predicted branching fraction versus the X^{++} mass.

Other effects of the GR scalar bosons are addressed briefly below:

(i) Neutral-current-to-charged-current ratio in neutrino scattering. Transitions of weak bosons to scalars give one-loop radiative corrections to W,Z masses. Using the formalism of Ref. 7, the deviation from unity of $\rho=\text{NC/CC}$ is

$$\Delta\rho=G_F M^2(1-\ln 2)/(\sqrt{2}\pi^2) .$$

At the one-standard-deviation level of the experimental value[8] $\rho=1.00\pm0.02$, the bound on the mass of X^{++} is $M<280$ GeV.

(ii) Anomalous magnetic moment of muon. The contribution[9] of GR scalar bosons to $a=(g-2)_\mu$ is

$$a=-5(g^2)_{\mu\mu}{m_\mu}^2/(48\pi^2M^2) .$$

The limit[10]

$$a_{\text{expr}}-a_{\text{QED}}>-18\times10^{-9}$$

on the difference of QED and experimental values

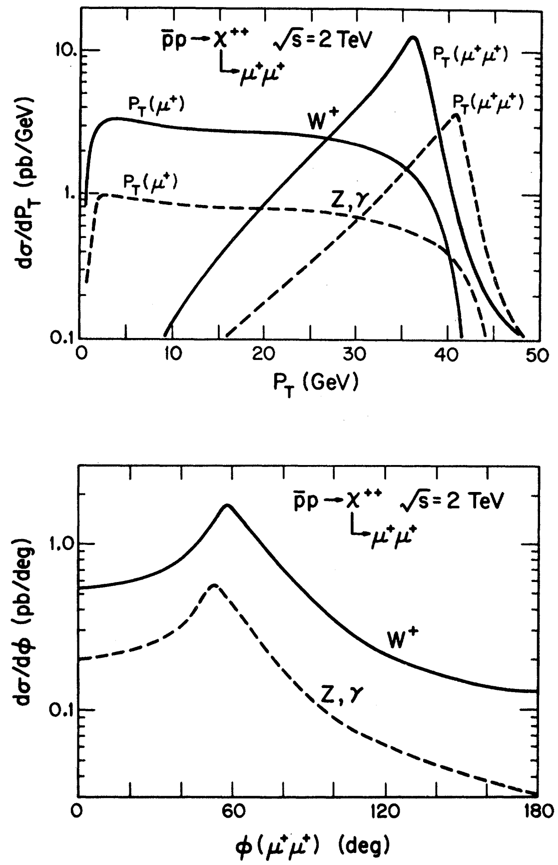

FIG. 3. Distributions in $\bar{p}p$ collisions at $\sqrt{s}=2$ TeV of muons from X^{++} production and leptonic decay. (a) Transverse momentum of μ^+ and $\mu^+\mu^+$. (b) Azimuthal angle between the two μ^+ in the transverse plane. The case $M=20$ GeV and $g_{ll'}=10^{-3}$ is illustrated.

only requires that $M>0.14$ GeV for $g_{ll'}=10^{-3}$, for all l,l'.

(iii) Muonium (μ^+e^-)—antimuonium (μ^-e^+) transition. The angular frequency of muonium-antimuonium transitions[11] due to the X^{++} is

$$\omega=g_{ee}g_{\mu\mu}(m_e\alpha^3)/(\pi M^2) .$$

For $g_{ee}=g_{\mu\mu}=10^{-3}$, the experimental limit[12] $\omega<83$ kHz only requires that $M>0.04$ GeV.

(iv) Narrow X^{--} resonance in $e^-e^-\to l^-l'^-$. The total width of X^{--} is

$$\Gamma(\text{keV})=0.4(M/20)\sum_{i,j}(g_{ij}/10^{-3})^2 + 0.022(M/20)^5 , \qquad (9)$$

where M is in GeV units. Taking all $g_{ll'}$ equal to 10^{-3}, the width is only 3.6 keV if M is 20 GeV. The integrated cross section over the X^{--} reso-

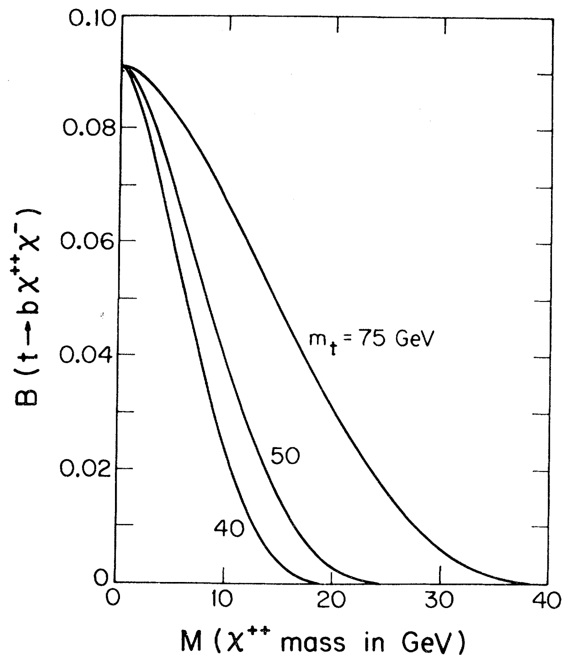

FIG. 4. Predicted branching fraction for the t-quark decay mode $t \to b\chi^{++}\chi^{-}$ versus the χ^{++} mass.

nance is

$$\int d\sqrt{s}\,\sigma(e^-e^- \to l^-l'^-) = 4\pi^2 B_{ee} B_{ll'} \Gamma / M^2 \,, \tag{10}$$

where

$$B_{ll'} = \Gamma(\chi^{--} \to l^- l'^-)/\Gamma \,.$$

For $M = 20$ GeV and all $g_{ll'} = 10^{-3}$, $\int d\sqrt{s}\,\sigma = 15$ MeV-nb when summed over all $l^- l'^-$ channels. The cross section is about 10 to 20 nb for 1 MeV energy resolution.

(v) Correction to QED *in* $e^+e^- \to \mu^+\mu^-$ *from* χ^{++} *exchange.* The modifications to cross section are orders of magnitude below the weak contribution and are thus unobservable.

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[1]G. B. Gelmini and M. Roncadelli, Phys. Lett. 99B, 411 (1981).

[2]H. M. Georgi, S. L. Glashow, and S. Nussinov, Nucl. Phys. B193, 297 (1981); S. L. Glashow, in *Second Workshop on Grand Unification, Ann Arbor, 1981,* edited by J. P. Leveille, L. R. Sulak, and D. G. Unger (Birkhauser, Boston, 1981).

[3]V. Barger, W. Y. Keung, and S. Pakvasa, Phys. Rev. D 25, 907 (1982).

[4]G. B. Gelmini, S. Nussinov, and M. Roncadelli, Report No. MPI-PAE/PTh 59/81 (unpublished).

[5]J. D. Vergados, Phys. Lett. 109B, 96 (1982).

[6]S. M. Korenchenko *et al.*, Zh. Eksp. Teor. Fiz. 70, 3 (1976) [Sov. Phys.—JETP 43, 1 (1976)]; K. G. Hayes and M. L. Perl, Report No. SLAC-PUB-2699, 1981 (unpublished).

[7]M. B. Einhorn, D. R. T. Jones, and M. Veltman, Nucl. Phys. B191, 146 (1981).

[8]J. E. Kim, P. Langacker, M. Levine, and H. H. Williams, Rev. Mod. Phys. 53, 211 (1981).

[9]J. Leveille, Nucl. Phys. B137, 63 (1978).

[10]J. Bailey *et al.*, Phys. Lett. 67B, 225 (1977); J. Calmet *et al.*, Rev. Mod. Phys. 49, 21 (1976).

[11]G. Feinberg and S. Weinberg, Phys. Rev. 123, 1439 (1961).

[12]D. Bechis *et al.*, University of Maryland Report No. ORO 2504-292, 1981 (unpublished).