Phenomenology of Massive Vectorlike Doublet Leptons

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Massive vectorlike electroweak doublets are generic in many extensions of the standard model. Even though one member of the doublet is necessarily electrically charged these particles are not easily detected in collider experiments. The neutral and charged states within the doublet are split by electroweak symmetry breaking. In the absence of mixing with other states, the radiatively generated splitting is in the range 200–350 MeV for $m \gtrsim \frac{1}{2} m_Z$. The charged state decays to the neutral one with an $\mathcal{O}$(cm) decay length, predominantly by emission of a soft charged pion. The best possibility to detect these massive charged particles is to trigger on hard initial state radiation and search for two associated soft charged pions with displaced vertices. The mass reach for this process at LEPII is limited by luminosity rather than kinematics. [S0031-9007(98)06501-6]

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The best probe for physics beyond the standard model is direct production of new states at high energy colliders. Many theories of electroweak symmetry breaking require additional states which are charged under electroweak gauge interactions. In many cases the states form chiral representations of $SU(2)_L \times U(1)_Y$, and necessarily gain mass only from electroweak symmetry breaking. However, it is possible for states with electroweak scale mass to transform under vector representations, even though there may be no symmetry apparent in the low energy theory which protects them from gaining a large mass. For example, the masses of vectorlike fermions can be protected by global chiral symmetries which are spontaneously broken at the electroweak scale. Likewise, in supersymmetric theories, matter supermultiplets which transform under a vector representation, and are massless at the high scale, remain massless to all orders due to the nonrenormalization theorem. The fermionic components of such supermultiplets can gain mass from a field which is a singlet under $SU(2)_L \times U(1)_Y$, but nonetheless gains an expectation value in association with electroweak symmetry breaking.

Massive vector representations can naturally carry a conserved or approximately conserved quantum number. This can forbid or highly suppress mixing with standard model fermions, and render the lightest state of the representation effectively stable on the scale of an accelerator experiment.

In this paper we discuss the phenomenology of a massive stable vector fermion doublet of $SU(2)_L$. In grand unified theories this representation arises in $\mathbf{5} \oplus \overline{\mathbf{5}} \in SU(5)$ or equivalently if a standard model generation is embedded in $\mathbf{27} \in E_6$ [1]. Vector representations of this type may also be required in theories of low scale supersymmetry breaking in which $U(1)_{PQ}$ and $U(1)_{R-PQ}$ Higgs sector symmetries are spontaneously broken at the electroweak scale [2]. Furthermore, the fermionic partners of the up- and down-type Higgs bosons in supersymmetry form a vectorlike $SU(2)_L$ doublet. These Higgsinos become mass eigenstates in the limit $|m_\lambda^2 - \mu^2| \gg m_Z^2$, where m_λ are the gaugino masses. The analysis given below becomes applicable in this limit if the Higgsino is the lightest supersymmetric particle.

Surprisingly, even though one member of the doublet is necessarily charged, these states turn out to be very difficult to detect experimentally. It is usually assumed that, if kinematically accessible, a heavy charged particle is easily detected at a high energy collider. This is generally true if (i) the heavy charged particle is nonrelativistic and lives long enough to pass through the entire detector, depositing a greater than minimum ionizing track, or (ii) it decays promptly to visible final states with energetic charged leptons and/or jets. The electrically charged state of the vectorlike $SU(2)_L$ doublet discussed here satisfies neither (i) nor (ii). The decay length to the neutral state, although macroscopic, is too short to allow direct triggering on the primary charged tracks. In addition, the visible decay products are too soft to allow direct triggering. However, as discussed below, triggering on associated initial state radiation allows a search for decays over a macroscopic distance to the very soft charged particles in the final state.

The representation considered here is a pair of $SU(2)_L$ doublet Weyl fermions with $U(1)_Y$ hypercharge $Y = \pm 1$, where $Q = T_3 + \frac{1}{2} Y$ [3]. This may be represented as a single Dirac fermion

$$L = \begin{pmatrix} L^0 \\ L^- \end{pmatrix}_{Y=-1}. \tag{1}$$

Other hypercharge assignments are not unifiable in a conventional manner, and have both members of the doublet charged. This representation is referred to as a doublet lepton since the left-handed component has the same gauge quantum numbers as a left-handed standard model lepton. This Dirac state can gain an $SU(2)_L \times U(1)_Y$ invariant mass, $\mathcal{L} \supset -m\bar{L}L = -m(\bar{L}^+ L^- + \bar{L}^0 L^0)$. In

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the absence of mixing with standard model leptons, the lowest order operator which can split L^- and L^0 in the presence of $SU(2)_L \times U(1)_Y$ breaking is $\bar{L}T^a L H^\dagger T^a H$, where H is the Higgs boson operator. For fermionic doublets this is a nonrenormalizable operator. In a renormalizable theory it receives finite calculable corrections. The mass splitting $\delta m \equiv m_{L^\pm} - m_{L^0}$ is therefore calculable within the low energy theory [4].

At lowest order the mass splitting comes from one-loop corrections with virtual photon and Z boson exchange to both the masses and wave functions. Virtual $W^\pm$ bosons do not contribute since the couplings to L^- and L^0 are identical. The one-loop mass splitting for on-shell states is

$$\delta m = \frac{\alpha}{2} m_Z f(m_L^2/m_Z^2), \tag{2}$$

where $f(r)$ is the loop function

$$f(r) = \frac{\sqrt{r}}{\pi} \int_0^1 dx[2 - x]\ln\left[1 + \frac{x}{r(1 - x)^2}\right]. \tag{3}$$

For $r \ll 1$, $f(r) \to 0$ and for $r \gg 1$, $f(r) \to 1$. The radiatively generated mass splitting is plotted in Fig. 1 for m_L in the range 50–100 GeV. The asymptotic value of the splitting for $m_L^2 \gg m_Z^2$ is $\delta m = \frac{1}{2}\alpha m_Z \simeq 355$ MeV. In this limit the mass renormalization is twice as large in magnitude and opposite in sign as compared with wave function renormalization.

The important features of the radiatively induced mass splitting can be understood in an effective field theory analysis. In the low energy theory below m_Z the L^- mass receives a linearly divergent contribution from the virtual photon loop. This divergence is cut off in the full theory by momenta above $\mathcal{O}(m_Z)$ for which electroweak symmetry is effectively restored. The splitting is therefore proportional to the electromagnetic fine structure constant times the Z boson mass. The linear divergence in momentum space corresponds in real space to the Coulomb self-energy of L^-. In the heavy field limit of $m_L^2 \gg m_Z^2$, the mass splitting (2) is precisely the difference between the Coulomb self-energies of L^- and L^0 due to the photon and Z boson classical electric fields [5]. In this interpretation it is clear that L^- is heavier than L^0, and that the splitting vanishes without electroweak symmetry breaking.

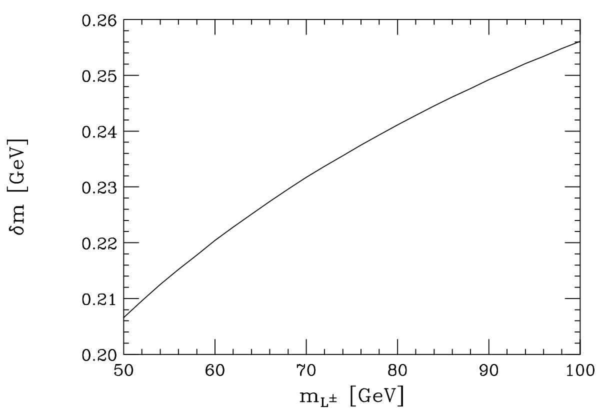

FIG. 1. Doublet mass splitting $\delta m \equiv m_{L^\pm} - m_{L^0}$ as a function of m_L.

The form of the splitting can also be understood in the effective theory above m_Z. In this description the coupling of the gauge eigenstates W^3 and B, of the $SU(2)_L$ and $U(1)_Y$ gauge groups, respectively, to L^- and L^0 are identical. Gauge invariance then implies that only diagrams which mix W^3 and B through an even number of Higgs insertions can contribute to the splitting. All of these effective operators receive infrared divergent contributions which are cut off by momenta of $\mathcal{O}(m_Z)$.

In a supersymmetric theory there are additional contributions to the mass splitting (2). At lowest order these come from one-loop diagrams with internal neutralinos and the scalar partner of the vector doublet. With an $SU(2)_L \times U(1)_Y$ invariant soft mass for the scalar partner of the form $\mathcal{L} = -m_{\tilde{L}}^2 \tilde{L}^\dagger \tilde{L}$, these contributions appear only as corrections to the vector doublet wave function. Electroweak symmetry breaking enters the supersymmetric loops at lowest order in two ways. The first is through gaugino-Higgsino mixing in the neutralino mass matrix. Since the lowest order operator in the effective theory above m_Z which splits L^0 and L^- requires at least two Higgs insertions, this contribution arises only at second order in gaugino-Higgsino mixing. In the mostly gaugino or Higgsino region of parameter space this contribution is then suppressed compared with (2) by $\mathcal{O}(m_L m_Z/(\mu^2 - m_\lambda^2))$. The second way electroweak symmetry enters is through the scalar partner $SU(2)_L$ D term. This splits the scalar $\tilde{L}^-$ and $\tilde{L}^0$ masses by $\mathcal{O}(m_Z^2/m_L)$. At one loop this modifies the vector doublet splitting by an amount which is suppressed compared with (2) by $\mathcal{O}(m_L m_Z/m_{\tilde{L}}^2)$. The loop momenta for both types of supersymmetric contributions are $\mathcal{O}(\max(m_{\tilde{L}}, m_\lambda))$. Because of these inherent suppressions, over much of the parameter space possible supersymmetric contributions are small compared with the dominant standard model contribution (2) to the doublet mass splitting. Small corrections are, however, sensitive to the superpartner spectrum.

The neutral state of the doublet, L^0, is rendered effectively stable by discrete or continuous global chiral symmetries. The state L^- can, however, decay to L^0 via charged current interactions. For the mass range of interest here the most important decay modes are $L^\pm \to L^0\pi^\pm$, $L^0 e^\pm \nu$, and $L^0 \mu^\pm \nu$. The partial widths for these modes are

$$\Gamma(L^\pm \to L^0\pi^\pm) = \frac{G_F^2}{\pi} \cos^2\theta_c f_\pi^2 \delta m^3 \sqrt{1 - b_\pi^2}, \tag{4}$$

$$\Gamma(L^\pm \to L^0 l^\pm \nu) = \frac{G_F^2}{15\pi^3} \delta m^5 \sqrt{1 - b_l^2}\, P(b_l), \tag{5}$$

where,

$$P(b_l) = 1 - \frac{9}{2} b_l^2 - 4 b_l^4 + \frac{15 b_l^4}{2\sqrt{1 - b_l^2}} \tanh^{-1} \sqrt{1 - b_l^2}\,, \qquad (6)$$

$f_\pi \simeq 130$ MeV, θ_c is the Cabibbo angle, $b_\pi = m_\pi/\delta m$, and $b_l = m_l/\delta m$. The branching ratios for each of these modes are plotted in Fig. 2 as a function of m_L. The exclusive mode $L^\pm \to L^0 \pi -$ of course dominates since it is two body and since $f_\pi \sim \delta m$.

Observable signatures of a massive vector doublet are very limited. Virtual contributions to the oblique electroweak parameters S and T are insignificant since the doublet does not gain a mass from electroweak symmetry breaking. The S parameter is proportional to corrections to mixing between W^3 and B gauge eigenstates. This requires at least two Higgs insertions and arises only at two loops. The T parameter is proportional to isospin violation, which likewise arises only at two loops. Direct decay of the Z boson to massive doublets is, however, important if kinematically open. The contribution to the Z boson total width is equivalent to $2[1 + (1 - 2\sin^2\theta_W)^2] \simeq 2.6$ massive Majorana neutrino species. This would unacceptably modify the Z width unless $m_L \gtrsim \frac{1}{2} m_Z$.

Direct detection of doublets which are too heavy to affect the Z boson total width is very challenging even though they are produced copiously if kinematically accessible; at an e^+e^- collider, $\sigma(L^+L^-) \sim \sigma(L^0\bar{L}^0) \sim \sigma(\mu^+\mu^-)\sqrt{1 - 4m_L^2/s}$ [6]. The neutral L^0 and $\bar{L}^0$ interact weakly like a massive neutrino and exit the detector without depositing visible energy. The principle reason for the difficulty in observing $L^\pm$ is that the decays (4) and (5) give an $L^\pm$ decay length of $\mathcal{O}$(cm). The laboratory frame $L^\pm$ decay distance for different center of mass energies relevant at the CERN Large Electron-Positron Collider (LEPII) is shown in Fig. 3. The typical decay length is unfortunately too short to utilize a topological trigger which identifies essentially back-to-back charged tracks largely independent of total energy deposition in the detector. Such a trigger requires that at least one of the tracks traverse the inner tracking region which typically extends to $\mathcal{O}$(30 cm). Triggering on the very soft charged decay products is equally difficult. For the $L^\pm \to L^0\pi^\pm$ decay mode, the $\pi^\pm$ radius of curvature in the detector magnetic field is $\mathcal{O}$(m/T). Separating such tracks at the trigger level from soft charged tracks arising from beam-beam interactions is problematic.

One method to search for production of invisible or nearly invisible particles is to trigger on an associated hard radiated photon. This has been suggested for counting neutrino species [7], and as a means to search for neutral supersymmetric particles, including photinos [8,9], neutralinos [10,11] sneutrinos [9,10,12], and nearly degenerate Higgsinos or W-inos [13]. In the approximation that the associated photon arises solely from initial state radiation, a photon radiator function [14] can be convoluted with the radiation-free cross section to obtain the differential cross as a function of $c_\gamma \equiv \cos\theta_\gamma$ and $x_\gamma = E_\gamma/E_{\rm beam}$:

$$\frac{d\sigma(L^+L^-\gamma)}{dx_\gamma dc_\gamma} = \sigma(L^+L^-)\,[(1 - x_\gamma)s]R(x_\gamma, c_\gamma; s)\,, \qquad (7)$$

where,

$$R(x_\gamma, c_\gamma; s) = \frac{\alpha}{\pi}\frac{1}{x_\gamma}\left[\frac{1 + (1 - x_\gamma)^2}{1 + 4m^2/s - c_\gamma^2} - \frac{x_\gamma^2}{2}\right]. \qquad (8)$$

The LEP experiments can trigger on central photons with $|\cos\theta_\gamma| \lesssim 0.7$ and energies greater than 5–10 GeV [15]. The cross section $\sigma(e^+e^- \to L^+L^-\gamma)$ at $\sqrt{s} = 183$ GeV with this photon coverage is plotted in Fig. 4 for several values of m_L as a function of the minimum photon energy for tagging, $E_\gamma^{\rm min}$.

The largest backgrounds for single hard photon production are $e^+e^- \to \nu\bar{\nu}\gamma$ and $e^+e^- \to Z\gamma$ with $Z \to \nu\bar{\nu}$,

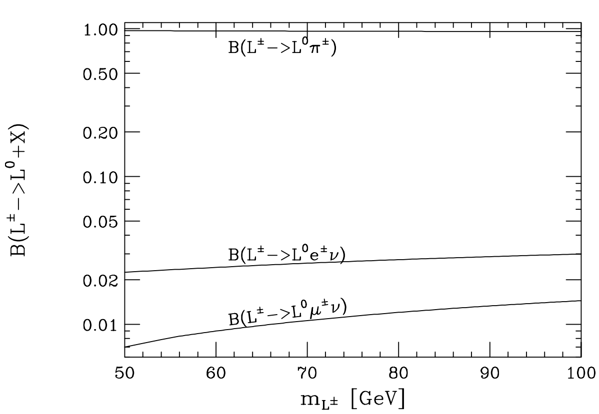

FIG. 2. Branching ratios for $L^\pm \to L^0 X$ where $X = \pi^\pm$, $e^\pm\nu$, or $\mu^\pm\nu$.

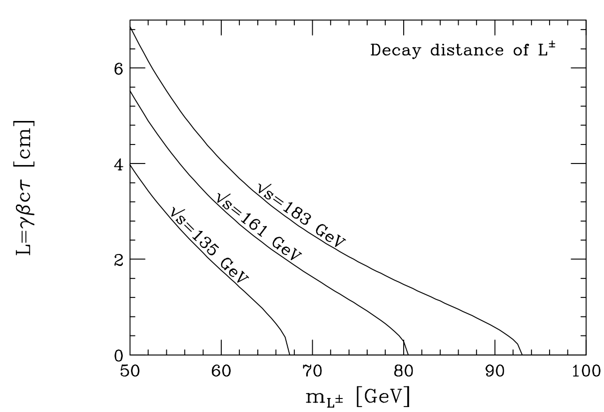

FIG. 3. Decay distance of $L^\pm$ in the laboratory frame at an e^+e^- collider as a function of m_L. The decays are boosted for $\sqrt{s} = 135$, 161, and 183 GeV.

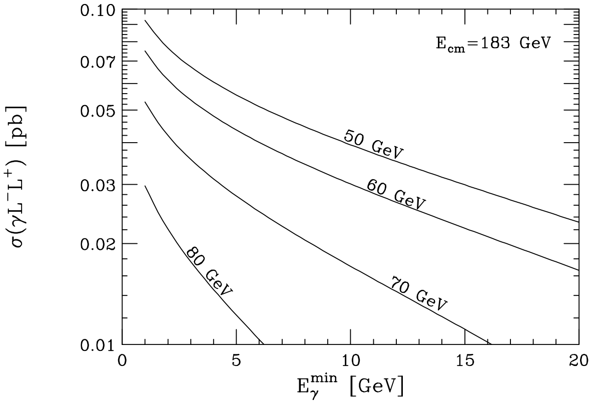

FIG. 4. Total cross section $\sigma(e^+e^- \to L^+L^-\gamma)$ at $\sqrt{s} = 183$ GeV with $|\cos\theta_\gamma| < 0.7$ and $E_\gamma > E_\gamma^{\min}$ as a function of $E_\gamma^{\min}$ for various values of m_L.

and to a lesser extent $e^+e^- \to l^+l^-\gamma$ with both l^+ and l^- forward and undetected. However, at the analysis level these can be separated from the signal $e^+e^- \to L^+L^-\gamma$ by requiring identification of two soft $\pi^\pm$ and/or $l^\mp$ arising from $L^\pm$ decays. Additional processes with very soft $\pi^\pm$ and/or $l^\mp$ in association with a photon may provide a small background. These can be separated from the signal by requiring that the soft tracks have some nonvanishing impact parameter with the beam axis. Pairs of soft charged tracks with displaced vertices in association with a hard photon provide a striking signature for massive vector doublets.

With an integrated luminosity of 240 pb^{-1} between the four LEP experiments at $\sqrt{s} = 183$ GeV, we estimate that the analysis described above could be sensitive to a doublet mass up to roughly 70 GeV. This assumes five signal events with $E_\gamma^{\min} = 8$ GeV and $|\cos\theta_\gamma| < 0.7$. A full Monte Carlo simulation of the signal and background with complete detector performance folded in would probably reduce this reach slightly.

Unlike many other signatures at e^+e^- machines, the experimental reach for vectorlike lepton doublets is not limited merely by the center of mass energy, but rather by the luminosity of the accelerator. Future LEPII runs with higher energy and much higher luminosity will greatly increase the search capability for these particles. Extension of searches to future lepton colliders such as the Next Linear Collider is also possible. Searches for L^+L^- production in association with a photon or Z boson are also in principle possible at hadron colliders. The larger background of soft charged tracks from beam-beam interactions, however, makes identification of displaced charged tracks from $L^\pm$ decay more challenging.

Finally, it is interesting to note that independent precision measurements of m_L from the total cross section, and of δm from the $L^\pm$ decay length distribution and decay pion energy spectrum, would be sensitive to deviations from the standard model one-loop mass splitting (2). This would provide an indirect probe for additional states beyond the photon and Z boson which can couple the heavy doublet to electroweak symmetry breaking through virtual processes, such as in supersymmetric theories.

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