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EFFECTS OF A NEW HEAVY NEUTRAL VECTOR BOSON AT SLC AND LEP*

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ABSTRACT

Models with an extra neutral gauge boson (Z') are discussed. We constrain the Z' mass as a function of its mixing angle with the known Z^0 by requiring that the Z^0 mass not be shifted excessively by this mixing, and from the Higgs vacuum expectation value structure of the mass matrix. We compare these limits with those previously found from neutral current experiments. We discuss possible effects of non-excluded models on e^+e^-physics at SLC and LEP.

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1. Introduction

The standard model, though experimentally successful, is not a fundamental theory. A recurring feature of attempted improvements involves embedding it in a larger symmetry group which is exact at higher energies. Many such theories, particularly grand unified theories and superstring theories, introduce new U(1) gauge symmetries which could remain unbroken to rather "low" energies. These extensions thus have extra flavor diagonal Z bosons, possibly with masses as low as 120 GeV.[1,2]

In this talk I will briefly describe three such models. I first describe the Lagrangian and discuss the relevant couplings. I then derive two kinds of limits on the mass of the new Z' versus its mixing angle (θ) with the familiar Z. The first limit comes from requiring that M_Z not be shifted excessively by mixing with the Z' (since the measured value of M_Z is in reasonable agreement ($\pm$ few GeV) with the standard model prediction). The second comes from imposing restrictions due to the Higgs sector of such models, which determines the $Z-Z'$ mass matrix. Finally, I discuss the effects the new bosons could have on polarization and forward-backward asymmetries.These effects may be seen at e^+e^-colliders such as SLC and LEP.

This talk is based on work done in collaboration with Fred Gilman.[3] Cvetič and Lynn,[4] and Bélanger and Godfrey,[5] have also considered effects due to extra Z''s at e^+e^- colliders.

2. Preliminaries

The neutral current interaction Lagrangian can be written as:

$$\mathcal{L}_{NC} = eA_\mu J^\mu_{em} + g_z Z_\mu J^\mu_Z + g' Z'_\mu J^\mu_{Z'} \tag{2.1}$$

where the couplings are given by

$$g_z = \frac{e}{\sqrt{x_W - x_W^2}} \qquad \text{and} \qquad g' = \frac{e}{\sqrt{1-x_W}} \quad . \tag{2.2}$$

Here J^μ_{em} is the electromagnetic current, $J^\mu_Z = J^\mu_{3L} - x_W Q^\mu$ is the standard $Z-$boson current, we define $J^\mu_{Z'} = \bar{f}_L\gamma^\mu\widetilde{Q}_L f_L + \bar{f}_R\gamma^\mu\widetilde{Q}_R f_R$ (the charge $\widetilde{Q}$ is a number dictated by the group structure), and $x_W = \sin^2\theta_W$. The amplitude for electron–positron annihilating to a fermion–antifermion pair is then

$$\mathcal{A}_{IJ} = \begin{pmatrix} g_{Z_0} & g_{Z'_0} \end{pmatrix}_{f_J} \begin{pmatrix} s - M^2_{Z_0} + iM_{Z_0}\Gamma_{Z_0} & \delta m^2 \\ \delta m^2 & s - M^2_{Z'_0} + iM_{Z'_0}\Gamma_{Z'_0} \end{pmatrix}^{-1} \begin{pmatrix} g_{Z_0} \\ g_{Z'_0} \end{pmatrix}_{e_I} + \mathcal{A}_\gamma . \tag{2.3}$$

Strictly speaking, we need to consider the Z' width to be energy dependent in order for the mass mixing approach to be exact. However, for the asymmetries and off-resonance cross-sections that I discuss in this talk, the width of the Z' has negligible effect.

In the expression for the amplitude, $(g_{Z_0})_{e_I}$ is the coupling of the Z_0 to appropriately-handed electrons (I stands for left or right); similarly for the Z_0' and for the fermions. Here the subscript 0 indicates the unmixed states. As an example, let us consider the couplings to left-handed electrons, which carry the charges $Q = -1$, $I_3 = -1/2$, $\tilde{Q}_\eta = -1/6$ (η denotes a particular model that I shall describe shortly). We have

$$\begin{aligned} (g_{Z_0})_{e_L} &= g_Z(I_3 - Qx_W) = \frac{3(-1+2x_W)}{2\sqrt{x_W - x_W^2}} = -.203 \\ (g_{Z_0'})_{e_L} &= g' \cdot \tilde{Q}_\eta = \frac{-.3}{6\sqrt{1-x_W}} = -.056 \end{aligned} \tag{2.4}$$

where we use the value $x_W = .22$. In general we have

$$\frac{g_{Z_0}}{g_{Z_0'}} = \frac{g_Z}{g'} \times \frac{I_3 - Qx_W}{\tilde{Q}} = 2.14 \times \frac{Qx_W - I_3}{\tilde{Q}} \tag{2.5}$$

for coupling to a fermion with charges Q, I_3 and $\tilde{Q}$.

Let us now consider some specific examples of E_6 breaking. I follow the notation of Ref. 2. The first is $E_6 \to SO(10) \times U(1)_\psi$; this particular breaking pattern has also been studied in Refs. 6 and 2. The 27 of E_6 contains all the fermions of one generation, and decomposes as follows:

$$(27)_{E_6} \to (16, x) + (10, -2x) + (1, 4x) \tag{2.6}$$

where the first number of each pair is the $SO(10)$ representation and the second is the $U(1)_\psi$ charge ($\tilde{Q}_\psi$); $x = -\frac{1}{6}\sqrt{\frac{5}{2}}\sqrt{\lambda_\psi} = -.26\sqrt{\lambda_\psi}$. The E_6 breaking scale determines λ_ψ, which is one if the initial E_6 breaking occurs at the same scale as the breaking to $SU(3)_c \times SU(2)_L \times U(1)$, and smaller if the $U(1)$ breaks off at higher energies. For definiteness, we take $\lambda_\psi = 1$. The $SO(10)$ multiplets break to $SU(5)$ multiplets in the usual way:

$$\begin{aligned} (16) &\to (10)[u, d, \bar{u}, e^+] + (5^*)[\bar{d}, \nu, e^-] + (1)[\bar{N}_L] \\ (10) &\to (5)[D, \bar{E}^0, E^+] + (5^*)[\bar{D}, E^0, E^-] \\ (1) &\to (1)[S_L^0] \quad . \end{aligned} \tag{2.7}$$

In this example, all members of one $SO(10)$ multiplet have the same $\tilde{Q}$ charge. The particles shown are all left-handed; the charges of right-handed fermions are given by $\tilde{Q}(f_R) = -\tilde{Q}(\bar{f}_L)$. For comparison I list the value of the charge $I_3 - Qx_W$ for all the particles of the 27:

e^-	e^+	u	$\bar{u}$	d	$\bar{d}$	ν	N, S_L^0	D	$\bar{D}$	E^0	E^-	E^+
$-.28$	$-.22$	.35	.15	$-.43$	$-.07$	.5	0	.07	$-.07$	.5	$-.28$	$-.22$

Note that while the couplings of the Z and Z' to normal fermions are comparable, the Z' has much larger couplings to the exotic fermions.

Another breaking, typical of some left-right symmetric theories, is $E_6 \to SO(10) \to SU(5) \times U(1)_\chi$, with the 27 decomposing as follows:

$$(27)_{E_6} \to (10, x) + (5^*, -3x) + (1, 5x) + (5, -2x) + (5^*, 2x) + (1, 0). \tag{2.8}$$

Here the first number of each pair is the $SU(5)$ representation and the second is the $U(1)_\chi$ charge; $x = \frac{1}{2\sqrt{6}}\sqrt{\lambda_\chi} = .204\sqrt{\lambda_\chi}$ and the first 5^* is the one that comes from the 16 of $SO(10)$. λ_χ ranges from 1 (if the breaking is at the $SU(5)$ scale) to 2/3 (if the breaking is at the Planck mass). Again, we choose $\lambda_\chi = 1$.

The third scheme that we consider has an extra $U(1)$ that is a linear combination of those of the first two:[7]

$$\widetilde{Q}_\eta = \sqrt{\frac{3}{8}}\,\widetilde{Q}_\chi - \sqrt{\frac{5}{8}}\,\widetilde{Q}_\psi \quad . \tag{2.9}$$

The fermion U(1) charges are given by the decomposition

$$(27)_{E_6} \to \left(10, \frac{1}{3}\right) + \left(5^*, -\frac{1}{6}\right) + \left(1, \frac{5}{6}\right) + \left(5, -\frac{2}{3}\right) + \left(5^*, -\frac{1}{6}\right) + \left(1, \frac{5}{6}\right) \quad . \tag{2.10}$$

Again, all these charges should be multiplied by $\sqrt{\lambda_\eta}$, a number of the order of unity (which we take to be unity for definiteness).

The decay width for Z' to some channel, say $\nu\bar{\nu}$, can be written down in terms of the Z width, scaled to the appropriate mass:

$$\Gamma(Z' \to \nu\bar{\nu}) = \left(\frac{g_{Z'}}{g_Z}\right)^2 \Gamma_{\text{at } M_{Z'}}(Z \to \nu\bar{\nu}) = \left(\sqrt{x_W}\frac{\widetilde{Q}(\nu)}{I_3(\nu)}\right)^2 \frac{g_Z^2}{4\sqrt{2}}\frac{M_{Z'}}{12\pi\sqrt{2}} = .00153 \cdot \widetilde{Q}^2(\nu) \cdot M_{Z'}. \tag{2.11}$$

We can express the decay width for Z' into any other channel as the ratio of the appropriate charges, squared, times the above width. For example,

$$\Gamma(Z' \to \text{all}) = \frac{\sum_{\text{all}} \widetilde{Q}^2}{\widetilde{Q}_\nu^2}\Gamma(Z' \to \nu\bar{\nu}). \tag{2.12}$$

The total widths of the Z' in the three models, and the branching ratios into known fermions, are:

	$\Gamma(Z' \to \text{all})$	$BR(e^+e^-)$	$BR(u\bar{u})$
ψ	4.8 (22.95)	6.1% (1.3%)	1.83% (.39%)
χ	10.5 (22.95)	4.4% (2.01%)	13.2% (6.03%)
η	5.75 (22.95)	3.7% (.93%)	17.7% (4.5%)

The numbers in parentheses correspond to the assumption that all exotics are light enough that the Z' can decay to them; the other numbers represent the other extreme—no exotics below threshold for Z' decay. The decay widths are given in units of $10^{-3} \times M'_Z$. The $d\bar{d}$ branching ratios are three times the electron branching ratios (the three arises from colour).

3. Limits on Masses Versus Mixing Angles

I begin by checking which theories are already ruled out by current limits on M_Z, as the mixing can significantly shift the Z mass. We can write the $Z - Z'$ mass matrix in the form

$$M^2 = M_{Z_0}^2 \begin{pmatrix} 1 & a \\ a & b \end{pmatrix} \Longrightarrow \begin{pmatrix} M_1^2 & 0 \\ 0 & M_2^2 \end{pmatrix}, \tag{3.1}$$

where the diagonalization is implemented by the unitary rotation $\begin{pmatrix} \cos\theta & -\sin\theta \\ \sin\theta & \cos\theta \end{pmatrix}$. We may relate the above parameters through the following expressions:

$$\tan 2\theta = \frac{2a}{1-b} \qquad M_1^2 = \frac{M_{Z_0'}^2 \sin^2\theta - M_{Z_0}^2 \cos^2\theta}{\sin^2\theta - \cos^2\theta} \approx M_{Z_0}^2 + (M_{Z_0}^2 - M_{Z_0'}^2)\theta^2 \tag{3.2}$$

The mass shift in the Z_0 due to mixing is given by

$$M_Z - M_{Z_0} = \frac{M_{Z_0}^2 - M_{Z_0'}^2}{2M_{Z_0}}\,\theta^2. \tag{3.3}$$

If we require that the Z be shifted by at most 3 GeV (from its value in the standard model, given a value of x_W) we have $|\theta| \leq .2$ for $M_Z \approx 150$ GeV; $|\theta| \leq .1$ for $M_Z \approx 250$ GeV. Figure 1 shows the limits on $M_{Z'}$ versus θ coming from limiting the Z shift to be at most 3 GeV; and at most 1 GeV. These bounds are model-independent in that they only depend on the charged gauge boson sector of the theory and not on the couplings or Higgs structure. In a model with no extra W's, such as the η model, favored by superstrings, we can accurately determine what we expect M_{Z_0} to be, using $M_{Z_0}^2 = M_W^2 / \cos^2\theta_W$. With present measurements of M_Z, shifts of about 3 GeV are excluded;* further measurements can vastly improve this limit (or discover a shift). In what follows, we require that the maximum shift be 3 GeV below 93 GeV. In a theory with extra W's as well as extra Z's, such as a left-right symmetric theory which involves the additional gauge group $SU(2)_R$, the W mass will also be shifted by mixing. Both masses will be shifted in the same direction; it is possible that their ratio may not change significantly. The observation of a shift in M_Z will indicate new physics, but a null observation may not place a good bound. In these models we will have to rely on a determination of $\sin^2\theta_W$ which does not depend on knowing M_Z and M_W, so that M_{Z_0} will be known to poorer accuracy, and limits on $\theta - M_{Z'}$ space will be weaker.

Another limit comes from the Higgs sector of the model. For example, for the $U(1)_\eta$ discussed above, the Z and Z' masses are generated by two Higgs doublets—H and H', and

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* We compare the measured value of M_Z with its value as calculated from the measured value of M_W, using $M_W = 38.65/\sin\theta_W$ (at the one-loop level) to fix $\sin\theta_W$. The UA1 experiment measures M_Z to be 1 GeV smaller than this "theoretical" value determined from M_W. Thus, taking the statistical errors of $M_Z(th)$ and $M_Z(exp)$ in quadrature, and ignoring the systematic errors (which come from an energy scale uncertainty which should affect both numbers approximately equally) we find that a 3 GeV decrease in M_Z from $M_Z(th)$ is about a 1.1σ effect. UA2, however, measures M_W to be smaller; $M_Z(th)$ is slightly smaller than $M_Z(exp)$, and a 3 GeV decrease from $M_Z(th)$ is a 2σ effect. We use numbers from Ref. 8.

one Higgs singlet—N. These particles have vacuum expectation values (VEV) v_1, v_2 and χ, respectively. The charges of these particles are H—$(1, \frac{4}{3})$, H'—$(-1, \frac{1}{3})$, and N—$(0, -\frac{5}{3})$ where I have given the Z charge, then the Z'. Recalling that $g'/g_Z = \sqrt{x_W}$, we have

$$M^2 = M_{Z_0}^2 \begin{pmatrix} 1 & \sqrt{x_W}\,\dfrac{4v_1^2 - v_2^2}{3(v_1^2 + v_2^2)} \\ \sqrt{x_W}\,\dfrac{4v_1^2 - v_2^2}{3(v_1^2 + v_2^2)} & x_W\,\dfrac{16v_1^2 + v_2^2 + 25\chi^2}{9(v_1^2 + v_2^2)} \end{pmatrix}, \tag{3.4}$$

using $M^2 = \sum Q_i^2 v_i^2$. The off diagonal element of this matrix can range from $-\frac{1}{3}\sqrt{x_W}$ to $\frac{4}{3}\sqrt{x_W}$, while $M_{Z'}$ is essentially free to vary independently. Thus we have the following bound on the mixing angle:

$$\frac{-\frac{8}{3}\sqrt{x_W}}{\frac{M_{Z'}^2}{M_Z^2} - 1} < \tan 2\theta < \frac{\frac{2}{3}\sqrt{x_W}}{\frac{M_{Z'}^2}{M_Z^2} - 1}. \tag{3.5}$$

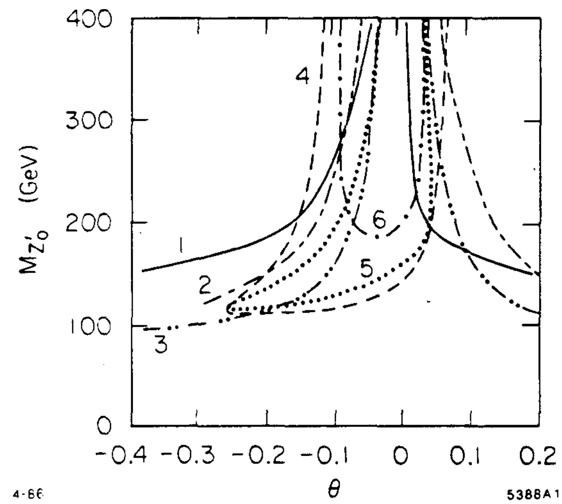

Fig. 1.

In Fig. 1, I show various limits on $\theta - Z'$ space. All lines are lower bounds. Line 1 is the limit due to the Higgs VEV structure of the mass matrix (for the η model—similar limits can be derived for the other models). Lines 2 and 3 are the bounds from constraining the Z mass shift to 3 and 1 GeV respectively. Lines 4,5 and 6 are from Durkin and Langacker,[2] who have considered the limits arising from neutral current experiments and the W and Z masses. Line 4 is for the η model, 5 for the η model with restrictions on the Higgs sector corresponding to what we have assumed, and 6 is for the χ model.

4. Effects on e^+e^- Physics

Figure 2 shows the polarization and forward-backward asymmetries in $e^+e^- \to \mu^+\mu^-$ for a model not excluded by other constraints: the η model with no mixing and $M_{Z'} = 150$ GeV. The asymmetries are given by

$$A_{Pol} = \frac{|\mathcal{A}_{RR}|^2 - |\mathcal{A}_{LL}|^2}{\sum_{IJ} |\mathcal{A}_{IJ}|^2} \qquad \text{and} \qquad A_{FB} = \frac{|\mathcal{A}_{RR}|^2 + |\mathcal{A}_{LL}|^2 - 2|\mathcal{A}_{LR}|^2}{\sum_{IJ} |\mathcal{A}_{IJ}|^2} \tag{4.1}$$

where $\mathcal{A}_{IJ}$ is the amplitude given in Eq. (2.3). The dotted line is the Z_0 alone, for comparison. Deviations from the standard model are present a little above the Z.

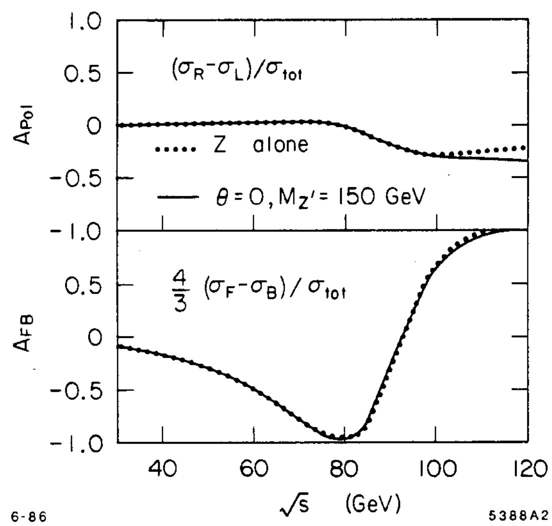

Fig. 2.

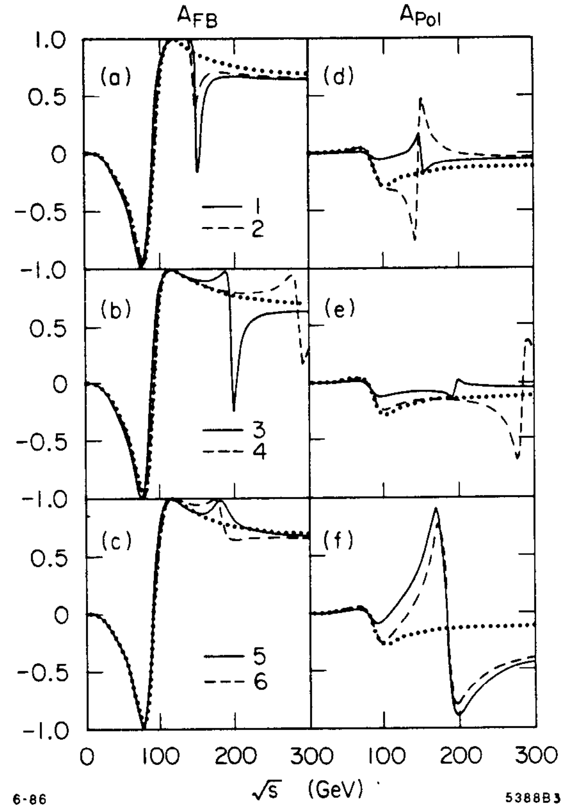

Fig. 3.

The Z' will be much easier to detect (or exclude) if it mixes with the Z. In Fig. 3 I show A_{Pol} and A_{FB} for a number of models. Figures 3(a),(d) show effects of the η model, for $M_{Z'} = 150$ GeV, and $\theta = 0$ (line 1) and $-.2$ (line 2). Figures 3(b),(e) show effects of the η model, for $(M_{Z'}, \theta) = (200, -.15)$ (line 3), and $(295, -.05)$ (line 4). Figures 3(c),(f) show effects of the χ model, for $(M_{Z'}, \theta) = (200, -.1)$ (line 5), and $(200, 0)$ (line 6). The ψ model gives similar, but smaller, effects. None of these models are excluded by present physics. The dotted line in all figures is the Z_0 alone for comparison. Effects in the quark channels are even more impressive, but will be harder to measure due to the difficulty of accurately predicting hadronization processes.

For a discussion of observable effects on cross-sections see Ref. 3.

5. Conclusions

We have looked at bounds on mixing angles versus extra boson mass in models with extra Abelian gauge symmetries. By requiring that the Z mass not be shifted excessively, which is a (somewhat) model-independent constraint, we get the bounds $|\theta| \leq .2$ for $M_Z \approx 150$ GeV; $|\theta| \leq .1$ for $M_Z \approx 250$ GeV. We can also get bounds from the Higgs VEV structure of the mass matrix, e.g, for the η model the approximation

$$.15 \left(\frac{M_Z}{M_{Z'}}\right)^2 > \theta > -.6 \left(\frac{M_Z}{M_{Z'}}\right)^2 \tag{5.1}$$

holds for large $M_{Z'}$. These limits are roughly comparable to the neutral current limits. Models that are not excluded by these limits could have prominent effects at SLC and LEP.

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