SEARCHING FOR A NEUTRAL E_6 GAUGE BOSON AT HERA

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In some superstring models, compactification leads to an E_6 gauge group which is further broken to an effective $SU(3)_c \times SU(2)_L \times U(1) \times U(1)$ gauge symmetry. The present experimental constraints on the neutral gauge boson associated with the extra U(1) group allow the existence of the second Z in the observability range of HERA. We study in detail the signals in the asymmetries measurable in polarized $e^{\mp}p$ collisions and point out the experimental precision required for detection.

The agreement of the $SU(3)_c \times SU(2)_L \times U(1)_Y$ gauge theory with all available experimental data is quite remarkable. Yet, this model contains too many unexplained structures and parameters to be a truly fundamental theory. In the hope to find answers to some of the open questions one has studied modifications and extensions in various directions such as compositeness, supersymmetrization and grand unification including gravity. Particularly great attention is being paid to the current development of superstring theories [1] because of their potential to lead to a "theory of everything" which, in the low-energy limit, approximates the so successful standard gauge model.

From the point of view of examining experimentally the idea that the familiar particles and interactions originate from a superstring in higher dimensions somewhere above the Planck scale, it is very important to study all imaginable possibilities of how the effective low-energy theory may deviate from the standard model. Popular examples [2] are the minimal $SU(3)_c \times SU(2)_L \times U(1) \times U(1)$ gauge theories resulting from the Wilson loop breaking of a supersymmetric E_6 grand-unified theory in the process of compactification of the heterotic $E_8 \times E_8$ superstring in ten dimensions. Although the appearance of one extra U(1) in the low-energy limit is by no means a theorem [3], it certainly represents an interesting case for phenomenological studies. The consequent deviations from the standard model must of course be rather small at present energies, nevertheless, they may become observable in the near future at one or the other of the new high-energy colliders.

In this letter, we investigate the effects of a heavy neutral gauge boson Z′ associated with a new U(1) symmetry in ep collisions at HERA. Since one can rotate the neutral gauge sector of $SU(2)_L \times U(1) \times U(1)$ such that the first U(1) coincides with the usual weak hypercharge group, the modified neutral current lagrangian in the fermion sector,

$$\mathscr{L} = eJ^{\mu}_{\mathrm{em}}A_{\mu} + \frac{e}{\sin\theta_{\mathrm{w}}\cos\theta_{\mathrm{w}}}J^{\mu}_{\mathrm{NC}}Z_{\mu} + g_{Y'}J'^{\mu}Z'_{\mu}, \tag{1}$$

differs formally only by the last term from the familiar lagrangian of the standard electroweak model. More definitely, in eq. (1), A_{μ} is the photon field coupled to the electromagnetic current

$$J^{\mu}_{\mathrm{em}} = \sum_f \bar{f}\gamma^{\mu}Q_f f, \tag{2}$$

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Q_f being the em charge. Z_μ is the ordinary Z field coupled to the standard neutral current

$$J^\mu_{\mathrm{NC}} = \sum_f \left[\bar{f}_{\mathrm{L}} \gamma^\mu \left(T_{3f} - \sin^2\theta_{\mathrm{w}}\, Q_f \right) f_{\mathrm{L}} + \bar{f}_{\mathrm{R}} \gamma^\mu \left(-\sin^2\theta_{\mathrm{w}}\, Q_f \right) f_{\mathrm{R}} \right], \tag{3}$$

T_{3f} being the third component of the weak isospin and $f_{\mathrm{L,R}} = \frac{1}{2}(1 \mp \gamma_5) f$ denoting the left(L)- and right(R)-handed components of the fermion fields. Here, the Weinberg angle θ_{w} is defined as in the standard model. Finally, Z'_μ is the new Z′ field coupled to the current

$$J'^\mu = \sum_f \left[\bar{f}_{\mathrm{L}} \gamma^\mu Y'_{f_{\mathrm{L}}} f_{\mathrm{L}} + \bar{f}_{\mathrm{R}} \gamma^\mu Y'_{f_{\mathrm{R}}} f_{\mathrm{R}} \right], \tag{4}$$

$Y'_{f_{\mathrm{L,R}}}$ being the new U(1) charge. Inspired by the superstring scenario sketched above we assume an underlying E_6 structure. In this case, the 15 helicity components of a standard fermion family form, together with new fermion species, a fundamental 27-plet of E_6. The charges $Y'_{f_{\mathrm{L,R}}}$ are then fixed by specifying the breaking pattern [3] of $E_6 \to SU(3)_c \times SU(2)_L \times U(1)_Y \times U(1)_{Y'}$. Here, we adopt a particular model considered in most of the preceding phenomenological investigations, e.g. in refs. [4–7], and leave the study of other possibilities to a subsequent publication. Since NC scattering at HERA only involves ordinary fermions (if one disregards the possibility of flavour changing neutral currents [8]), it is sufficient to give Y' for the latter. The assignment is

$$Y'_{\mathrm{u_L,d_L}} = -Y'_{\mathrm{u_R,e_R}} = 2Y'_{\mathrm{d_R}} = -2Y'_{\nu_{\mathrm{L}},\mathrm{e_L}} = 1/\sqrt{15} \tag{5}$$

for the lightest family and, correspondingly, also for the heavier families. Furthermore, the $U(1)_{Y'}$ gauge coupling has the value

$$g_{Y'} = \sqrt{5/3}\; e/\cos\theta_{\mathrm{w}}. \tag{6}$$

The normalization of eqs. (5) and (6) is such that in the E_6 symmetry limit (when $\sin^2\theta_{\mathrm{w}} = 3/8$) $g_{Y'}$ is equal to the $SU(2)_L$ coupling $g = e/\sin\theta_{\mathrm{w}}$.

In general, the Z and Z′ fields appearing in eq. (1) are not the physical fields which acquire definite masses through spontaneous breakdown of $SU(2)_L \times U(1)_Y \times U(1)_{Y'}$ to $U(1)_{\mathrm{em}}$. The mass eigenstates Z_1 and Z_2 are rather mixtures of Z and Z′,

$$\begin{pmatrix} Z_1 \\ Z_2 \end{pmatrix} = \begin{pmatrix} \cos\theta & \sin\theta \\ -\sin\theta & \cos\theta \end{pmatrix} \begin{pmatrix} Z \\ Z' \end{pmatrix}, \tag{7}$$

where, without assumptions on the Higgs sector, the mixing angle θ can take any value in the range $-\frac{1}{2}\pi \leqslant \theta \leqslant \frac{1}{2}\pi$. The state Z_1 with the lower mass eigenvalue m_{Z_1} is identified with the already observed neutral weak boson, whereas the heavier state Z_2 acquiring the mass m_{Z_2} is the object of our main interest. In the absence of mixing, m_{Z_1} obviously coincides with the mass m_Z of the standard Z-boson as can be seen from the relation

$$m_Z^2 = \cos^2\theta\, m_{Z_1}^2 + \sin^2\theta\, m_{Z_2}^2 \tag{8}$$

implied by eq. (7). This limit is also approached for very large values of m_{Z_2} as

$$\tan^2\theta = \left(m_Z^2 - m_{Z_1}^2 \right) / \left(m_{Z_2}^2 - m_Z^2 \right) \tag{9}$$

vanishes if $m_{Z_2} \to \infty$. In what follows we shall assume that the Higgs fields responsible for the spontaneous symmetry breaking are all $SU(2)_L$ doublets and singlets. Then [4],

$$\rho = m_W^2 / m_Z^2 \cos^2\theta_{\mathrm{w}} = 1 \tag{10}$$

where m_W is the charged weak boson mass and m_Z is given by eq. (8). This restriction again is motivated by the superstring scenario and supported [6] by the existing NC data. Rewriting now eq. (1) in terms of the physical bosons Z_1 and Z_2 one obtains

$$\mathscr{L} = eJ^{\mu}_{em}A_{\mu} + \frac{e}{\sin\theta_w \cos\theta_w}\left(J^{\mu}_1 Z_{1\mu} + J^{\mu}_2 Z_{2\mu}\right), \tag{11}$$

where

$$J^{\mu}_1 = \cos\theta\, J^{\mu}_{NC} + \sin\theta\, \sqrt{5/3}\, \sin\theta_w\, J'^{\mu}, \quad J^{\mu}_2 = -\sin\theta\, J^{\mu}_{NC} + \cos\theta\, \sqrt{5/3}\, \sin\theta_w\, J'^{\mu}, \tag{12}$$

with the currents J^{μ}_{NC} and J'^{μ} defined in eqs. (3) and (4), respectively. Thus, apart from m_{Z_1} and $\sin^2\theta_w$, the model introduced above contains only two more parameters: the mass m_{Z_2} of the new boson Z_2 and the mixing angle θ. These parameters are related by eq. (10) and are to be determined or constrained by experiment.

Present-day NC data put lower bounds on m_{Z_2} in the range from 100 to 150 GeV with θ not exceeding a few degrees [4–6]. As an example, the bound on m_{Z_2} and θ obtained by Durkin and Langacker [6] is shown in fig. 1. The possible existence of a second neutral boson with such a small mass makes the case very favourable for HERA, at least at first sight. However, it turns out that the effects of Z_2 on the NC cross sections are rather marginal. Even for m_{Z_2} as low as 100 GeV and at $Q^2 \gtrsim 10^4$ GeV^2, the highest momentum transfer accessible at HERA with reasonable statistics to perform an inclusive search, the cross sections do not deviate from the standard model predictions by more than a few percent. In order to discover the new boson or to improve the already existing bounds by a considerable amount, one must therefore study more sensitive quantities. Similarly as in other searches for new physics in ep collisions [9], the most sensitive experiments are asymmetry measurements. This fact can be exploited if longitudinally polarized $e^{\pm}$ beams are available as expected for HERA. As far as the attainable experimental precision is concerned, realistic estimates are left to the experts. We shall rather demonstrate the precision required in order to reach certain values of m_{Z_2} and θ.

Using $Q^2 = -q^2$ and $x = -q^2/2pq$ as independent variables, where q is the four-momentum of the exchanged boson and p is the incoming proton momentum, one can express the cross sections of deep inelastic $e^-_{L,R}p$ scattering calculated from eq. (11) as follows:

$$\frac{d\sigma(e^-_{L,R})}{dx\, dQ^2} = \frac{2\pi\alpha^2}{xQ^4}\left[1 + (1-y)^2\right]\left[F_2^{L,R}(x, Q^2) + h(y)xF_3^{L,R}(x, Q^2)\right]. \tag{13}$$

The variable y relates x and Q^2 to the total CM energy $\sqrt{s}$, $xys = Q^2$, and $h(y) = [1-(1-y)^2]/[1+(1$

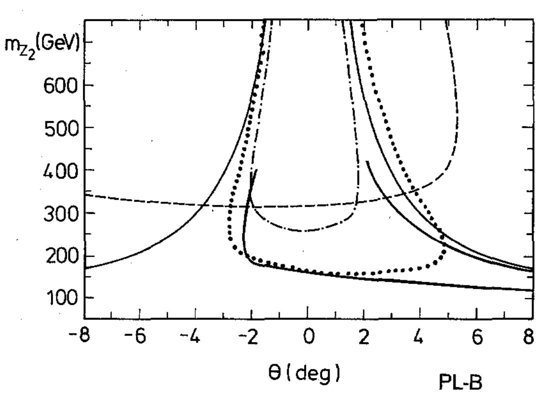

Fig. 1. Bounds on the mass and mixing angle of the second Z boson. Values inside the boundary drawn as thick full curves (from ref. [6] after changing the sign of θ) are compatible with present neutral current phenomenology. Values outside the thin full curves (from ref. [4]) are excluded by the observed W and Z masses. The other curves exhibit boundaries which could be obtained from measurements at HERA of the asymmetries A^{-+}_{RL} (dashed), A^{--}_{LR} (dashed–dotted) and A^{-+}_{RR} (dotted) with a precision $\delta A = 0.04$.

$-y)^2]$. Furthermore, the structure functions $F_2^{\mathrm{L,R}}$ and $xF_3^{\mathrm{L,R}}$ are given by

$$F_2^{\mathrm{L,R}}(x, Q^2) = \sum_f \left[xq_f(x, Q^2) + x\bar{q}_f(x, Q^2) \right] \tilde{F}_{2f}^{\mathrm{L,R}}(Q^2),$$
$$xF_3^{\mathrm{L,R}}(x, Q^2) = \sum_f \left[xq_f(x, Q^2) - x\bar{q}_f(x, Q^2) \right] \tilde{F}_{3f}^{\mathrm{L,R}}(Q^2), \tag{14}$$

where $q_f(x, Q^2)$ and $\bar{q}_f(x, Q^2)$ are the scaling violating quark and antiquark densities of flavor f inside the proton, respectively. The electroweak charges and propagators enter in the functions $\tilde{F}_{if}^{\mathrm{L,R}}(Q^2)$:

$$\tilde{F}_{2f}^{\mathrm{L,R}}(Q^2) = Q_f^2 + \sum_{i=1}^{2} \left[(v_{ie} \pm a_{ie})^2 (v_{if}^2 + a_{if}^2) P_i^2 - 2Q_f (v_{ie} \pm a_{ie}) v_{if} P_i \right]$$
$$+ 2(v_{1e} \pm a_{1e})(v_{2e} \pm a_{2e})(v_{1f} v_{2f} + a_{1f} a_{2f}) P_1 P_2,$$
$$\tilde{F}_{3f}^{\mathrm{L,R}}(Q^2) = \pm 2 \left(\sum_{i=1}^{2} \left[(v_{ie} \pm a_{ie})^2 v_{if} a_{if} P_i^2 - Q_f (v_{ie} \pm a_{ie}) a_{if} P_i \right] \right.$$
$$\left. + (v_{1e} \pm a_{1e})(v_{2e} \pm a_{2e})(v_{1f} a_{2f} + a_{1f} v_{2f}) P_1 P_2 \right). \tag{15}$$

The vector and axial vector charges v_{if} and a_{if} can readily be inferred from eqs. (3)–(5) and (12):

$$\begin{pmatrix} v_{1f} \\ v_{2f} \end{pmatrix} = \frac{1}{\sin 2\theta_{\mathrm{w}}} \begin{pmatrix} \cos\theta & \sin\theta \\ -\sin\theta & \cos\theta \end{pmatrix} \begin{pmatrix} T_{3f} - 2Q_f \sin^2\theta_{\mathrm{w}} \\ \sqrt{5/3}\, \sin\theta_{\mathrm{w}} (Y'_{f_{\mathrm{L}}} + Y'_{f_{\mathrm{R}}}) \end{pmatrix},$$
$$\begin{pmatrix} a_{1f} \\ a_{2f} \end{pmatrix} = \frac{1}{\sin 2\theta_{\mathrm{w}}} \begin{pmatrix} \cos\theta & \sin\theta \\ -\sin\theta & \cos\theta \end{pmatrix} \begin{pmatrix} T_{3f} \\ \sqrt{5/3}\, \sin\theta_{\mathrm{w}} (Y'_{f_{\mathrm{L}}} - Y'_{f_{\mathrm{R}}}) \end{pmatrix}, \tag{16}$$

while

$$P_i = Q^2 / (Q^2 + m_{Z_i}^2). \tag{17}$$

The positron cross sections $\mathrm{d}\sigma(\mathrm{e}^+_{\mathrm{L,R}})/\mathrm{d}x\, \mathrm{d}Q^2$ follow then directly from eq. (13) by replacing $F_2^{\mathrm{L,R}}$ by $F_2^{\mathrm{R,L}}$ and $xF_3^{\mathrm{L,R}}$ by $-xF_3^{\mathrm{R,L}}$. With the above results it is now straightforward to derive the desired asymmetries. In the short-hand $\mathrm{d}\sigma \equiv \mathrm{d}\sigma/\mathrm{d}x\, \mathrm{d}Q^2$, the polarization and charge asymmetries read

$$A_{\mathrm{LR}}^{\mp\mp} \equiv \frac{\mathrm{d}\sigma(\mathrm{e}_{\mathrm{L}}^{\mp}) - \mathrm{d}\sigma(\mathrm{e}_{\mathrm{R}}^{\mp})}{\mathrm{d}\sigma(\mathrm{e}_{\mathrm{L}}^{\mp}) + \mathrm{d}\sigma(\mathrm{e}_{\mathrm{R}}^{\mp})} = \frac{\pm(F_2^{\mathrm{L}} - F_2^{\mathrm{R}}) + h(y)(xF_3^{\mathrm{L}} - xF_3^{\mathrm{R}})}{(F_2^{\mathrm{L}} + F_2^{\mathrm{R}}) \pm h(y)(xF_3^{\mathrm{L}} + xF_3^{\mathrm{R}})} \tag{18}$$

and

$$A_{\substack{\mathrm{LL}\\ \mathrm{RR}}}^{-+} \equiv \frac{\mathrm{d}\sigma\left(\mathrm{e}^-_{\substack{\mathrm{L}\\ \mathrm{R}}}\right) - \mathrm{d}\sigma\left(\mathrm{e}^+_{\substack{\mathrm{L}\\ \mathrm{R}}}\right)}{\mathrm{d}\sigma\left(\mathrm{e}^-_{\substack{\mathrm{L}\\ \mathrm{R}}}\right) + \mathrm{d}\sigma\left(\mathrm{e}^+_{\substack{\mathrm{L}\\ \mathrm{R}}}\right)} = \frac{\pm(F_2^{\mathrm{L}} - F_2^{\mathrm{R}}) + h(y)(xF_3^{\mathrm{L}} + xF_3^{\mathrm{R}})}{(F_2^{\mathrm{L}} + F_2^{\mathrm{R}}) \pm h(y)(xF_3^{\mathrm{L}} - xF_3^{\mathrm{R}})}, \tag{19}$$

respectively, while the mixed asymmetries are given by

$$A_{\substack{\mathrm{LR}\\ \mathrm{RL}}}^{-+} \equiv \frac{\mathrm{d}\sigma\left(\mathrm{e}^-_{\substack{\mathrm{L}\\ \mathrm{R}}}\right) - \mathrm{d}\sigma\left(\mathrm{e}^+_{\substack{\mathrm{R}\\ \mathrm{L}}}\right)}{\mathrm{d}\sigma\left(\mathrm{e}^-_{\substack{\mathrm{L}\\ \mathrm{R}}}\right) + \mathrm{d}\sigma\left(\mathrm{e}^+_{\substack{\mathrm{R}\\ \mathrm{L}}}\right)} = h(y) \frac{xF_3^{\substack{\mathrm{L}\\ \mathrm{R}}}}{F_2^{\substack{\mathrm{L}\\ \mathrm{R}}}}. \tag{20}$$

Evidently, since one has only four different structure functions, eq. (14), only four of these asymmetries are independent. Nevertheless, it is useful to consider all six asymmetries and to compare their sensitivity to the presence of Z_2.

A first study of this kind has been performed by Angelopoulos et al. [7]. Our work extends this analysis in several respects. Firstly, the mixed asymmetries $A^{-+}_{LR,RL}$ are not considered in ref. [7]. As we shall show, A^{-+}_{RL} can provide the most sensitive test of the presence of Z_2 beyond the present limits. Secondly, whereas the above authors concentrate on a rather restricted range of values of m_{Z_2} and θ favored by a no-scale supergravity model [4], we analyze a considerably wider region of the parameter space. Thirdly, we also pay attention to the dependence of the results on the strategy followed in the analysis. This dependence is not negligible as substantiated later.

In eqs. (18)–(20) one can distinguish two effects arising from the presence of a second neutral boson: direct contributions to the asymmetries from Z_2-exchange and indirect modifications of the standard NC terms due to Z–Z′ mixing. Because of mixing the couplings of the lower mass eigenstate Z_1 to the ordinary fermions deviate from the standard NC couplings as detailed in eq. (16). Furthermore, m_{Z_1} differs from the standard model value m_Z according to eq. (8) and, consequently, the Weinberg angle obtained from the W and Z_1 mass ratio,

$$\sin^2\tilde{\theta}_w = 1 - m_W^2/m_{Z_1}^2, \tag{21}$$

is different from the angle defined by the $SU(2)_L$ and $U(1)_Y$ gauge couplings,

$$\sin^2\theta_w = g_Y^2/(g^2 + g_Y^2). \tag{22}$$

Using $g \sin\theta_w = g_Y \cos\theta_w = e$, $g^2 = 4\sqrt{2}\, G_F m_W^2$ and eqs. (8) and (10) one can rewrite eq. (22) as follows:

$$\sin^2\theta_w = (\mu/m_W)^2 = \tfrac{1}{2}\left(1 - \sqrt{1 - 4\mu^2/m_Z^2}\right) \tag{23}$$

with $\mu = 38.65$ GeV. The last number includes the standard loop corrections [10]. Small modifications in the present model due to the second boson and, possibly, new E_6 fermions (and superpartners) are neglected. Considering m_{Z_2} and θ as essentially free parameters one may now follow different strategies in fixing the other two parameters, m_{Z_1}, and $\sin^2\theta_w$, of the model.

(I) One starts at a given value of $\sin^2\theta_w$ (for the later illustrations we take $\sin^2\theta_w = 0.22$), calculates m_Z from eq. (23) and determines m_{Z_1} as a function of m_{Z_2} and θ from eq. (8).

(II) One starts at a given value of m_{Z_1} (for the later illustrations we take $m_{Z_1} = 93.3$ GeV), calculates m_Z from eq. (8) and determines $\sin^2\theta_w$ as a function of m_{Z_2} and θ from eq. (23).

Since the Weinberg angle is still more accurately known than the Z mass, i.e., m_{Z_1}, strategy (I) is appropriate for the present. However, by the time of the HERA experiments the Z mass will be known very precisely from measurements at SLC and LEP. Then strategy (II) is more appropriate. Of course, the mixing induced shifts of m_{Z_1} in (I) and $\sin^2\theta_w$ in (II) from the standard model values are (and will be even more strongly) restricted by experiment. These constraints can directly be turned into bounds on m_{Z_2} and θ. For instance, requiring [4] $\sin^2\theta_w - \sin^2\tilde{\theta}_w \leqslant 0.035$ in accordance with CERN $p\bar{p}$ collider data [11] and following strategy (II) one obtains the contour plotted in fig. 1. A similar curve results from approach (I) if $m_{Z_1} \geqslant 91$ GeV is imposed. We emphasize that these bounds solely reflect mixing effects and, hence, do not constrain m_{Z_2} for sufficiently small θ. In the latter region, the main effects come from interactions involving directly Z_2. Except at resonance, these effects are suppressed by the Z_2 propagator and, consequently, only visible if m_{Z_2} is sufficiently light. The above remarks are nicely illustrated in fig. 1 by comparing the bound [4] derived from the shift in $\sin^2\theta_w$ with the contour obtained in ref. [6] which includes the constraints from all low-energy NC processes. As a final remark, for a given process the size of the effects induced by Z–Z′ mixing and their interplay with the direct Z_2 contributions may differ quite

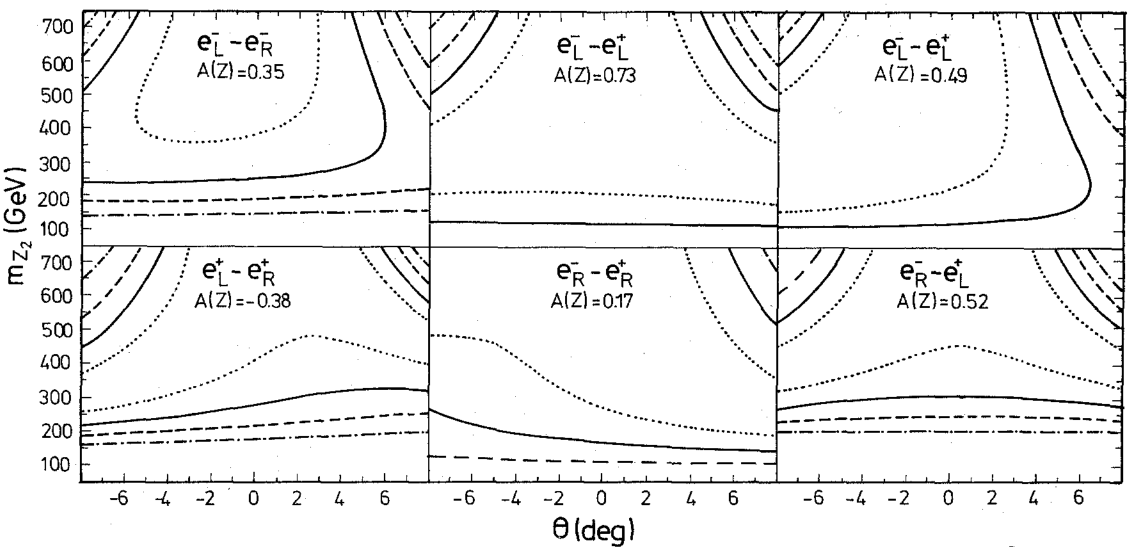

Fig. 2. Contours showing the values of m_{Z_2} and θ for which the asymmetries $A(\mathrm{Z})$ and $A(\mathrm{Z}_1, \mathrm{Z}_2)$ predicted in the standard model and in the two-Z model, respectively, differ by $\delta A \equiv |A(\mathrm{Z}_1, \mathrm{Z}_2) - A(\mathrm{Z})| = 0.02$ (dotted curves), 0.04 (full curves), 0.06 (dashed curves) and 0.08 (dashed–dotted curves). Here, strategy (I) described in the text is used in the analysis. The various asymmetries are characterized by the incoming lepton states while the numbers denote the respective standard model values.

drastically in the approaches (I) and (II). This is, in particular, the case for the ep asymmetries under consideration.

We now proceed to present our numerical results. To this end, we define

$$\delta A = |A(\mathrm{Z}) - A(\mathrm{Z}_1, \mathrm{Z}_2)| \tag{24}$$

for any of the asymmetries given in eqs. (18)–(20). Here, $A(\mathrm{Z})$ and $A(\mathrm{Z}_1, \mathrm{Z}_2)$ denote the predictions of the standard and the two-Z model, respectively. The contours in the (m_{Z_2}, θ)-plane corresponding to $\delta A = 0.02, 0.04, 0.06, 0.08$ are shown in fig. 2 for strategy (I) and in fig. 3 for strategy (II). As a somewhat extreme but not inconceivable kinematical point we have chosen $x = 0.3$ and $Q^2 = 2 \times 10^4$ GeV^2 at the HERA CM energy $\sqrt{s} = 314$ GeV. The absolute values of the various asymmetries expected in the standard model for $m_Z = 93.3$ GeV and $\sin^2\theta_w = 0.22$ are also given in the figures. In all calculations we have used set I of the Duke and Owens quark distribution functions [12]. Several comments are in order.

(i) In all cases one can clearly distinguish the regions in m_{Z_2} and θ where the effects are mainly due to Z_2-exchange (flat parts of the contours around $\theta \simeq 0$) as opposed to the regions where the effects dominantly originate in Z–Z′ mixing (steep part of the contours at large values of m_{Z_2}).

(ii) With few exceptions, the asymmetries react quite uniformly on the mixing induced effects at large values of m_{Z_2}. In contrast, for lower Z_2 masses one observes characteristic differences in the behaviour reflecting the model-dependent interferences of Z_2 with Z_1 and the photon. In other words, only for a relatively light Z_2 can one hope to discriminate its existence from other interpretations of a possible signal.

(iii) The indirect mixing effects are more pronounced in approach (II) than in approach (I). This is expected since the asymmetries are more sensitive to $\sin^2\theta_w$ [fixed in (I)] than to m_{Z_1} [fixed in (II)]. Obviously, for $\theta = 0$ both strategies lead to the same result as can be checked in figs. 2 and 3.

(iv) The mixed asymmetry A_{RL}^{-+} is the most sensitive probe of the presence of a second boson at moderate masses. On the other hand, in the case of a very heavy Z_2, the strongest constraints on the Z–Z′ mixing result from A_{LR}^{--} in approach (II).

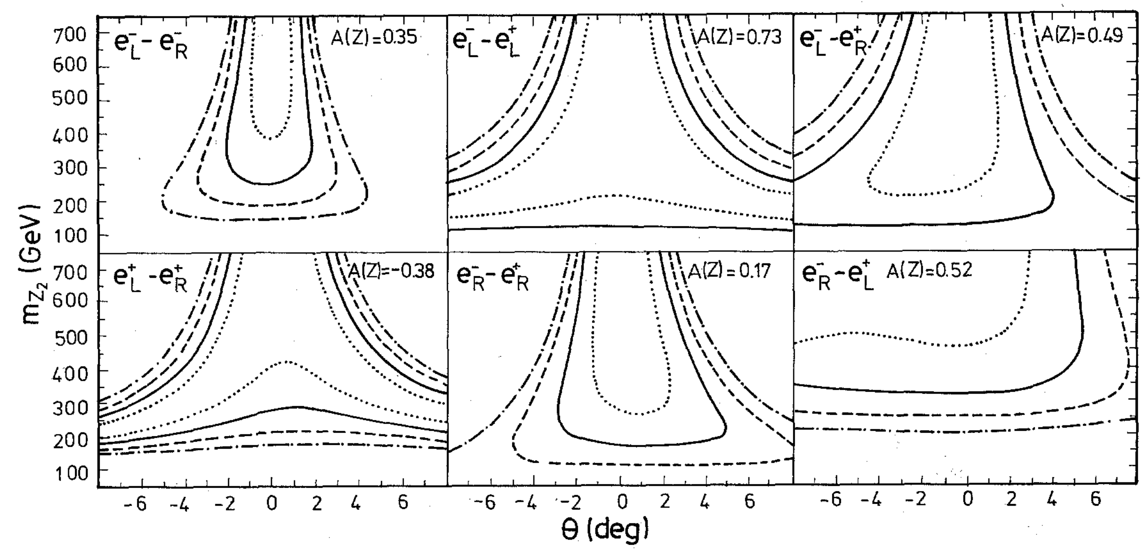

Fig. 3. Same as fig. 2, however, following strategy (II) described in the text.

As discussed earlier, fig. 1 shows the regions in the (m_{Z_2}, θ)-plane which are already excluded by present-day experiments. For clarity of the drawings, these bounds have not been included in figs. 2 and 3. In order to facilitate an assessment of the improvement which can be expected from HERA, some cases of fig. 3 are reproduced in fig. 1. For this comparison, we have assumed the observability of an absolute deviation from the standard model asymmetries by 4%, i.e., $\delta A = 0.04$, and selected the more sensitive asymmetry of the polarization, charge and mixed type. Such an accurate search would improve the present bounds quite appreciably, in particular the lower limit on m_{Z_2}, or find an effect. From fig. 1 one deduces the detection limit

$$m_{Z_2} \gtrsim 320 \text{ GeV} \quad \text{and} \quad |\theta| \lesssim 2°. \tag{25}$$

A similar lower bound on m_{Z_2} is obtained from A_{RL}^{-+} in approach (I), fig. 2. Although the limit on θ is uninteresting in this case, the constraint on m_{Z_2} combined with the present bounds exhibited in fig. 1 essentially implies the same overall boundary in the (m_{Z_2}, θ)-plane as the one quoted in eq. (25). Quite generally, given the present constraints on Z_2 the mixed asymmetry A_{RL}^{-+} is the most favourable quantity for inclusive tests at HERA.

Using the results reported in this paper one can readily examine the reach in m_{Z_2} and θ for all asymmetries and also different experimental conditions. For example, if the attainable precision is worse than $\delta A = 0.06$ (0.08) the detection limit of Z_2 drops to $m_{Z_2} \simeq 250$ (200) GeV. Likewise if the precision exemplified in figs. 2 and 3 can be accomplished but only at lower values of Q^2, the sensitivity to the presence of a second Z decreases rapidly. For instance, at $x = 0.15$ and $Q^2 = 10^4$ GeV^2, the boundary eq. (25) roughly corresponds to $\delta A = 0.02$, whereas for $\delta A \gtrsim 0.04$ one is only sensitive to masses below $m_{Z_2} \simeq 200$ GeV. Finally, without polarized $e^\pm$ beams one must search for effects in the cross sections themselves or in the unpolarized charge asymmetry. Unfortunately, for all values of m_{Z_2} and θ considered in figs. 1–3 and, in particular, for the values compatible with present NC phenomenology, the effects appear to be too small to be detected.

To conclude, although it will not be easy at HERA to probe the existence of a second neutral weak boson with properties as expected in an E_6 grand-unified model motivated by the heterotic $E_8 \times E_8$ superstring, the issue is certainly worth every experimental effort. Further work is also assigned to theory, in particular, the complete computation of the radiative corrections [13] to the ep cross sections and asymmetries.

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