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New heavy gauge bosons in pp and $p\bar{p}$ collisions

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Experimental signatures are analyzed for the production of heavy gauge bosons beyond W and Z at pp and $p\bar{p}$ colliders, including the Fermilab Tevatron and proposed multi-TeV machines. Bosons include right-handed W's and various Z's (including the one expected if I_{3R} and $B-L$ are gauged separately, not just in the combination $Y_W=2I_{3R}+B-L$). Signatures include characteristic decay asymmetries, which can occur for both pp and $p\bar{p}$ reactions, and neutral heavy leptons in the final states.

I. INTRODUCTION

For several years, low-energy neutral- and charged-current weak interactions have pointed toward an effective $SU(2)\times U(1)$ theory[1] of the electroweak interactions.[2] This conclusion is strengthened by the discovery of the W and Z at their predicted masses.[3] In the SU(5) model which unifies this electroweak theory with color SU(3),[4] no new gauge bosons are expected up to masses of $10^{14}-10^{15}$ GeV. Only indirect experiments, such as the search for proton decay, are capable of searches at such high masses.

There are reasons to adopt a more empirical approach toward the possibility of new gauge bosons observable at accelerator energies. The SU(5) theory does predict the weak mixing angle correctly,[5] but proton decay has not yet been observed at the expected rate. Other unified theories involving groups larger than SU(5) can accommodate present bounds on the nucleon lifetime without difficulty, though the prediction of the weak mixing angle is lost. At the same time, these unified theories entail new gauge bosons which can be quite light.

In the present article, we discuss some prospects for observing and identifying the properties of new gauge bosons in the multi-TeV range. We are stimulated in part by the possibility that proton-proton and/or proton-antiproton collisions at center-of-mass energies up to 40 TeV may be attainable before the end of this century. A general overview of the physics at such energies has been given in Ref. 6.

We shall be concerned with several possibilities for new gauge bosons, which illustrate the variety of experimental signatures that could arise at multi-TeV energies. Our intention thus is not so much to break new theoretical ground as to aid in planning of detectors. This, in turn, can affect parameters of accelerator design.

The bosons we shall consider include (1) the right-handed analog W_R of the standard W; (2) a second Z boson, called Z_χ, which couples differently to different SU(5) representations and arises most naturally in SO(10); (3) a third Z boson, called Z_ψ, which couples universally to all known fermions and arises naturally in E_6; and (4) neutral bosons with Yang-Mills structure, such as those which change generations or connect ordinary fermions with exotic ones.

Let us set the scale of possible masses and properties involved. We begin with the right-handed W.

The properties of a right-handed W were discussed phenomenologically by Bég, Budny, Mohapatra, and Sirlin,[7] though suggestions of the existence of such particles date back to some of the earliest attempts at grand unification.[8,9] In muon decays the absence of deviations from expected spectra were then used to conclude $M_{W_R}\gtrsim 3M_{W_L}$. Recent experiments have strengthened this bound somewhat.[10] This conclusion is based on the assumption that a light right-handed neutrino exists. In the absence of such a neutrino, no such conclusion can be drawn.

A much stronger limit can be set by examining the contribution of right-handed W's to the K_L-K_S mass difference.[11] The ordinary box diagram acquires a large enhancement when one of the W's is right-handed, leading to a bound

$$M_{W_R} > 1-2\ \text{TeV} \tag{1.1}$$

if the contribution of this box graph is not to exceed the observed value. While this value is somewhat dependent on the details of hadronic physics at long distances, it is sufficiently large that the right-handed W, if it exists, is unlikely to mix much with the left-handed W, and probably can be discussed in isolation to a large extent. We shall adopt this point of view in our discussion.

Further (somewhat model-dependent) limits on right-handed W masses have been placed with the help of nonleptonic decays,[12] the process $K\to\pi e\nu$,[13] and the neutron electric dipole moment.[14] It should be noted that not all authors[15] agree with the bound (1.1), but we shall assume it for purposes of discussion.

Heavy neutral gauge bosons (besides the standard Z) can manifest themselves in low-energy neutral-current phenomena such as those discussed in Refs. 2 and 16. There have been many discussions of the possible ways in which this can occur.[8,9,17–28]

A rule of thumb for heavy neutral gauge bosons, more or less independent of the specific models in which they arise, is that they will lead to modifications of magnitude $(M_{Z(\text{heavy})}/M_{Z_0})^2$ to observable quantities. Such quantities are typically measured to an accuracy of no more than 10%. This is true, for example, of the weak asymmetry in $e^++e^-\to\mu^++\mu^-$, or of the magnitude of parity violation in atomic physics. The lower bound on the corresponding heavy Z is then expected to be only about three times that of the standard Z. In practice the bounds are even weaker, so that low-energy phenomenology often can exclude heavy Z's only up to about 200 GeV in mass. (For one example of such a permissive model, see Ref. 27.) A model discussed earlier,[24] for which the major constraint comes from parity violation in atoms, has been analyzed by Bouchiat,[26] with the result that only Z's below about 150 GeV can be excluded.

Even for modestly heavy extra Z's, the effects of mixing with the light Z on the mass of this particle can be minimal. The examples mentioned above can lead to a lowering of the Z_0 mass from its predicted value by less than 2%. In what follows we shall ignore the effects of such mixings.

The remainder of this article is organized as follows. In Sec. II we discuss constraints on mixings of gauge bosons. This is done in order to avoid discussing such mixings in the remainder of the paper, which is rather concerned with direct searches for the new bosons. In such searches, if the bosons are heavy enough, their mixings with the standard W and Z are found to be unimportant.

Section III is devoted to couplings and mass scales of the new bosons. Section IV contains estimates of rapidity distributions and total cross sections, while Sec. V is concerned with forward-backward asymmetries in leptonic decays. Experimental signatures distinguishing the new bosons from the old ones (in respects aside from their masses) are mentioned in Sec. VI. Conclusions occupy Sec. VII.

II. MIXING EFFECTS

In extended electroweak models such as $\mathrm{SU}(2)\times\mathrm{U}(1)\times G$ one must consider the possibility of mixing between the W and Z and the new bosons in G. Fortunately, if the new bosons are heavy (e.g., $\gtrsim 1$ TeV) then one can show rather generally[29] that such mixings are severely limited by the closeness of the observed W and Z masses to the $\mathrm{SU}(2)\times\mathrm{U}(1)$ predictions.

For example, if G has a single neutral gauge boson Z_G then (after removing the photon) the neutral boson (mass)^2 matrix is

$$M^2=\begin{bmatrix} m_{Z_0}{}^2 & b \\ b & m_{Z_G}{}^2 \end{bmatrix}, \tag{2.1}$$

where $m_{Z_0}{}^2$ and $m_{Z_G}{}^2$ are the squares of the Z and Z_G masses in the absence of mixing and B is an arbitrary mixing (mass)^2. The physical mass eigenstates are then

$$Z_1=Z\cos\theta+Z_G\sin\theta$$

and

$$Z_2=-Z\sin\theta+Z_G\cos\theta\ .$$

It is easy to show for $m_{Z_0}\le m_{Z_G}$ that $m_{Z_1}{}^2<m_{Z_0}{}^2$, $m_{Z_2}{}^2>m_{Z_G}{}^2$ and that

$$\tan^2\theta=\frac{m_{Z_0}{}^2-m_{Z_1}{}^2}{m_{Z_2}{}^2-m_{Z_0}{}^2}\ . \tag{2.2}$$

In order to apply (2.2) it is necessary to assume that m_{Z_0} is given by the canonical $\mathrm{SU}(2)\times\mathrm{U}(1)$ value $93.8^{+2.4}_{-2.2}$ GeV.[2] That will be the case if the only Higgs fields with significant vacuum expectation values are SU(2) doublets (and singlets). Then, the observed values[3]

$$m_{Z_1}=\begin{cases} 95.6\pm1.4\pm2.9\ \text{GeV} & (\text{UA1}) \\ 91.9\pm1.3\pm1.4\ \text{GeV} & (\text{UA2}) \end{cases} \tag{2.3}$$

imply $\tan^2\theta=(0.073)^2=5.3\times10^{-3}$ for $m_{Z_2}=500$ GeV, $(0.036)^2=1.3\times10^{-3}$ for $m_{Z_2}=1$ TeV, and $(0.018)^2=3.3\times10^{-4}$ for $m_{Z_2}=2$ TeV.

Similar results hold for the mixing between the W and a single new charged boson. If there are two or more new neutral (or charged) bosons then the right-hand side of (2.2) becomes an upper bound on the squares of the relevant mixing angles.[29] Clearly, the agreement of the observed W and Z masses with the $\mathrm{SU}(2)\times\mathrm{U}(1)$ predictions strongly suggests that one can, to an excellent approximation, ignore the mixings of the heavy new bosons to the W and Z when discussing their production and decay. The only theories for which this would not hold would be models in which a significant increase in m_{Z_0} (from Higgs triplets, etc.) is fortuitously compensated by mixing effects. Such a possibility seems to us sufficiently unlikely that we will not consider it further.

III. COUPLINGS OF HEAVY W's AND Z's

In the present section we discuss couplings and branching ratios of several types of hypothetical particles a heavy, right-handed W, denoted W_R; various heavy Z's; and neutral gauge bosons with Yang-Mills couplings to one another.

A. Right-handed W

The couplings of left- or right-handed fermions to W_L (the standard W) or W_R may be expressed in terms of the Lagrangians

$$\mathscr{L}_{L,R}=-g_{L,R}\bar{\psi}_{L,R}\vec{\mathrm{I}}_{L,R}\cdot\vec{\mathrm{W}}_{L,R}\psi_{L,R}\ , \tag{3.1}$$

where

$$\psi_{L,R}=(1\mp\gamma_5)/2 \tag{3.2}$$

and

$$W^{\pm}_{L,R}=(W^1_{L,R}\mp iW^2_{L,R})/\sqrt{2}\ . \tag{3.3}$$

The coupling constant g_L is related to the Fermi constant G_F by

$$g_L{}^2/8M_W{}^2=G_F/\sqrt{2}\ , \tag{3.4}$$

so that it may be considered as known now that the W mass has been measured.[3]

The weak left-handed isospin multiplets are[30]

$$I_L=\tfrac{1}{2}:\quad \begin{bmatrix} u \\ d \end{bmatrix}_L,\ \begin{bmatrix} \nu_e \\ e^- \end{bmatrix}_L,\ \ldots, \tag{3.5}$$

$$I_L=0:\quad u_R,d_R,e_R^-\ ,\ \ldots\ . \tag{3.6}$$

The branching ratios of ordinary W_L's are then

$$\begin{aligned}&\Gamma(W_L^-\to e^-\bar{\nu}_e):\Gamma(W_L^-\to\mu^-\bar{\nu}_\mu):\Gamma(W_L^-\to\tau^-\bar{\nu}_\tau)\\ &:\Gamma(W_L^-\to d\bar{u}):\Gamma(W_L^-\to s\bar{c}):\Gamma(W_L^-\to b\bar{t})\\ &\qquad =1:1:1:3:3:3\ ,\end{aligned} \tag{3.7}$$

if the t quark is light enough that kinematic suppression is unimportant. Then

$$B(W_L^-\to e^-\bar{\nu}_e)=\tfrac{1}{12}=8.3\%\ . \tag{3.8}$$

What is different for right-handed W's?

The neutrino is known to exist in left-handed form, as $\nu_{iL}(i=e,\mu,\tau)$. The corresponding antineutrinos $\bar{\nu}_{iR}$ also exist, of course, and are also isodoublet members. There are several possibilities for the "right-handed neutrinos" ν_{iR} and their antiparticles $\bar{\nu}_{iL}$, assuming they exist.

(1) Neutrinos could be Dirac particles. In this case Eqs. (3.5)–(3.8) would be valid if one interchanged L and R everywhere.

(2) The ordinary neutrinos ν_{iL} could be Majorana particles, distinct (in mass) from ν_{iR}. In this case we shall denote

$$\nu_{iR}\equiv N_{iR}\ , \tag{3.9}$$

and the right-handed isospin multiplets are

$$I_R=\tfrac{1}{2}:\quad \begin{bmatrix} u \\ d \end{bmatrix}_R,\ \begin{bmatrix} N_e \\ e^- \end{bmatrix}_R,\ \ldots, \tag{3.10}$$

$$I_R=0:\quad u_L,d_L,e_L^-,\ \ldots\ . \tag{3.11}$$

The right-handed W's decay in a manner very similar to Eq. (3.7) except that ν_i is replaced by N_i. Thus, one signature for decay of W_R would be the production of (e.g.) N_e with branching ratio

$$B(W_R^-\to e^-\bar{N}_e)=\tfrac{1}{12}(M_{N_e}\ll M_{W_R})\ . \tag{3.12}$$

The subsequent behavior of N_e is model-dependent. It will be discussed further in Sec. VI. This new neutral lepton could decay or be stable, and its decay could take place immediately, after a detectable path length, or outside the detector. This wide range of possibilities means that W_R detection may be more experimentally challenging than that of the W_L, which used the distinctive signature (3.8).

The coupling constant g_R governs the rate of production of W_R in hadron-hadron collisions. We shall assume that at momentum transfers large compared with $M_{W_{L,R}}$, one has $g_L=g_R$. Such a relation occurs in many grand unified theories, as when SO(10) breaks down to SO(6)$\times$SO(4). The SO(4) is just $SU(2)_L\times SU(2)_R$, with equal coupling constants.

The asymmetry between $g_L(M_{W_L})$ and $g_R(M_{W_R})$ may be gauged from the relation (based on the renormalization group)

$$\frac{4\pi}{g_R{}^2}=\frac{4\pi}{g_L{}^2}+\frac{11-2N}{3\pi}\ln\frac{M_{W_R}}{M_{W_L}}\ . \tag{3.13}$$

We take the number of generations $N=3$, and $4\pi/g_L{}^2=29.2$ [for $G_F=1.16\times10^{-5}$ GeV^{-2}, $M_W\equiv M_{W_L}=81$ GeV in Eq. (3.4)]. Then for $M_{W_R}\le10$ TeV, we find

$$4\pi/g_R{}^2-4\pi/g_L{}^2\le2.6\ , \tag{3.14}$$

less than a 10% difference. For the rest of our discussion, we shall take

$$g_R=g_L \tag{3.15}$$

neglecting the small logarithmic variation implied by (3.13). (The value of $M_{W_R}\cong10$ TeV will turn out to be the highest that we can attain with colliding-beam machines which might be constructed in the foreseeable future.)

The W_R and W_L can mix with one another. In our simplified discussion, we shall neglect this mixing, as we expect it to be small for $M_{W_R}\gg M_{W_L}$.

B. Second Z in SO(10): Z_χ

The group SO(10) is an appealing candidate for unification of electroweak and strong interactions.[31] It has several advantages over the group SU(5):[4]

(1) In SO(10), each fermion generation belongs to a single 16-dimensional representation, composed of 5^*+10+1 of SU(5).

(2) The operations C and P are much simpler in SO(10):[32] $C(16)\to16$; $P(16)\to16^*$. In SU(5) C and P mix representations with one another in a rather obscure fashion.

(3) The observed value of the weak mixing angle, $\sin^2\theta=\frac{1}{5}-\frac{1}{4}$, can be accommodated in SO(10) without entailing a prediction of too rapid a proton decay.[33] Of course, in SO(10) the very precise SU(5) predictions of $\sin^2\theta$, in accord with present data, do not necessarily hold.

Since SO(10) has rank one higher than SU(5), it has one more neutral boson. We shall call this boson Z_χ. If Z_χ is fairly light, it can mix with the "standard" Z_0, with many interesting effects. The effects of a second Z on low-energy phenomena have been analyzed in detail.[17–28] Here we shall be concerned instead with direct production of the Z_χ in the limit $M_{Z_\chi}\gg M_{Z_0}$, when mixing can be neglected. The Z_χ then has very specific ratios of couplings to quarks and leptons, which we now analyze.

The Z_χ couples to a very specific U(1) charge which we shall call χ. This is the conserved charge of the group U(1) which arises when

$$\mathrm{SO}(10)\to\mathrm{SU}(5)\times\mathrm{U}(1)_\chi \,. \tag{3.16}$$

Each SU(5) representation in an SO(10) representation has a distinct value of χ.

We shall normalize χ in the same way as more familiar generators of SO(10). For a 16-plet, for example, $\sum I_{3L}{}^2=2$. We shall thus demand $\sum\chi^2=2$ as well for a 16-plet. Then we find, for members of a fermion generation,[24]

$$5^*\!: \; \chi=\frac{3}{2\sqrt{10}} \,, \; (\bar{d},e^-,\nu_e)_L \,, \tag{3.17a}$$

$$10\!: \; \chi=\frac{-1}{2\sqrt{10}} \,, \; (d,u,\bar{u},e^+)_L \,, \tag{3.17b}$$

$$1\!: \; \chi=\frac{-5}{2\sqrt{10}} \,, \; (\bar{N}_e)_L \,. \tag{3.17c}$$

Note that $5\chi(5^*)+10\chi(10)+\chi(1)$ vanishes, as it must, and that $\chi(1)-\chi(10)=\chi(10)-\chi(5^*)$. These two relations and the normalization $\sum\chi^2=2$ (for a 16-plet) form a convenient mnemonic for χ.

We then find

$$\begin{aligned}\Gamma(Z_\chi\to u\bar{u})&:\Gamma(Z_\chi\to d\bar{d}):\Gamma(Z_\chi\to e^+e^-)\\ &:\Gamma(Z_\chi\to\nu_e\bar{\nu}_e):\Gamma(Z_\chi\to N_e\bar{N}_e)\\ &=3(1^2+1^2):3(3^2+1^2):3^2+1^2:3^2:5^2\\ &=6:30:10:9:25 \,,\end{aligned} \tag{3.18}$$

$$B(Z_\chi\to u\bar{u})=\tfrac{1}{3}\times\tfrac{6}{80}=2.5\% \,, \tag{3.19a}$$

$$B(Z_\chi\to d\bar{d})=\tfrac{1}{3}\times\tfrac{30}{80}=12.5\% \,, \tag{3.19b}$$

$$B(Z_\chi\to e^+e^-)=\tfrac{1}{3}\times\tfrac{10}{80}=4.2\% \,, \tag{3.19c}$$

$$B(Z_\chi\to \nu_e\bar{\nu}_e)=\tfrac{1}{3}\times\tfrac{9}{80}=3.8\% \,, \tag{3.19d}$$

$$B(Z_\chi\to N_e\bar{N}_e)=\tfrac{1}{3}\times\tfrac{25}{80}=10.4\% \,, \tag{3.19e}$$

for three generations of quarks and leptons. Here we have assumed $M_N \ll M_{Z_\chi}/2$. If Z_χ cannot decay to $N\bar{N}$, its branching ratios to other fermion pairs will be raised accordingly.

The coupling of Z_χ to $u\bar{u}$ pairs is quite small. This is because both helicities of u quarks belong to the 10-plet, with small χ charge. This will make the Z_χ harder to produce in collisions for which valence quarks play a crucial role (since the valence quarks in a proton are predominantly u quarks, especially at high x). However, it could lead to some distinctive signatures in Z_χ decay, which we shall discuss in Sec. VI.

The branching ratio of Z_χ to e^+e^- is expected to be respectable (4.2%, as compared to about 3% for the Z_0). Charged-lepton pairs thus ought to be an acceptable way to detect Z_χ.

If $Z_\chi\to N_i\bar{N}_i$ is kinematically allowed, such decays form a major part of Z_χ final states. How they would show up depends on the properties of N_i, which we treat in Sec. VI.

The magnitude of the coupling of Z_χ to the χ charge is model-dependent. Several possibilities have been discussed in Ref. 24. Here we shall concentrate on one, to illustrate the method of calculation. Production cross sections can easily be rescaled as desired. As we shall show, the expected range of variation is not large, probably less than a factor of 2 in g^2.

We assume that the breaking (3.16) of SO(10) occurs at the same mass at which SU(5) breaks further to SU(3)×SU(2)×U(1) [the ability of SO(10) to resolve any problems with regard to proton decay is lost, but the successful SU(5) prediction of $\sin^2\theta$ is retained]. The coupling associated with χ (call it g_χ) then should have the same value at any Q^2 as that associated with weak hypercharge (call it g_Y) if the weak hypercharge is also normalized to $\sum Y^2=2$ for a 16-plet.

The value of g_Y is related to g_L and θ. To see this, we note that $g_Y\mathcal{N}=g'=e/\cos\theta$, where $\mathcal{N}$ is a normalization factor, and $g_L=e/\sin\theta$. Then

$$g_Y{}^2=g_L{}^2\tan^2\theta/\mathcal{N}^2 \,. \tag{3.20}$$

At the unification mass, $g_Y=g_L$ and $\tan^2\theta=\frac{3}{5}$,[4] so $\mathcal{N}^2=\frac{3}{5}$. Thus,

$$g_\chi=g_Y=\sqrt{5/3}\,g_L\tan\theta \,. \tag{3.21}$$

This is the value we shall assume at the Q^2 value appropriate to production of Z_χ. We shall neglect the variation of θ with Q up to M_{Z_χ} and simply take $\sin^2\theta=0.22$ from low-energy experiments.

If the breaking (3.16) of SO(10) were to occur at the Planck mass, $g_\chi{}^2$ would be reduced to about $\frac{2}{3}$ the value implied by Eq. (3.21).[24]

An "effective" Z_χ arises under much weaker assumptions than those given above. The electromagnetic charge may be written[8,22]

$$Q=I_{3L}+I_{3R}+\frac{B-L}{2} \,. \tag{3.22}$$

It couples to the photon. The Z_0 couples to an orthogonal linear combination of I_{3L} and $Y_W/2=I_{3R}+(B-L)/2$. If *any* neutral boson besides γ and Z_0 couples to any of the three quantities in (3.22), and if its mixing with γ and Z_0 can be neglected by virtue of its large mass, it must couple to the only linear combination of I_{3L}, I_{3R}, and $(B-L)/2$ orthogonal to both I_{3L} and Y_W. The definition of "orthogonal" requires that we normalize generators suitably.

If one is to gauge I_{3R} and $B-L$ separately, one must guarantee that they are anomaly-free. The fermions belonging to Eqs. (3.17a) and (3.17b) [5^*+10 of SU(5)] are free of anomalies in I_{3L} and $Y_W/2=I_{3R}+(B-L)/2$. However, they are not anomaly-free in I_{3R} and $B-L$ separately. The addition of the right-handed neutrino N is sufficient to banish anomalies in I_{3R} and $B-L$. With fermions belonging to all of Eqs. (3.17) [5^*+10+1 of SU(5)] one has

$$\sum I_{3L}{}^3=\sum I_{3R}{}^3=\sum(B-L)^3=0 \,, \tag{3.23}$$

and this is sufficient to remove all anomalies. But now we automatically have the normalization

$$\sum I_{3L}{}^2=\sum I_{3R}{}^2=2 \; . \tag{3.24}$$

If a generator proportional to $B-L$ is to be normalized similarly,

$$\sum a^2 \left[\frac{B-L}{2}\right]^2=2 \; , \tag{3.25}$$

where it turns out that $a^2=\frac{3}{2}$, there is then a combination of I_{3R} and $a(B-L)/2$ orthogonal to that in the weak hypercharge Y_W. It is just

$$\chi\equiv\frac{1}{\sqrt{10}}[2I_{3R}-\tfrac{3}{2}(B-L)] \; , \tag{3.26}$$

where we have chosen $\sum\chi^2=2$ as well for a set (3.17).

For convenience one may eliminate $B-L$ from Eqs. (3.22) and (3.26) and write

$$\chi=\frac{1}{\sqrt{10}}[5I_{3R}+3(I_{3L}-Q)] \; . \tag{3.27}$$

The values of I_{3L} and I_{3R} may be read off directly from Eqs. (3.5), (3.6), (3.10), and (3.11).

The Lagrangian for the coupling of Z_χ to fermions is simply

$$\mathscr{L}_{Z_\chi}=-g_\chi\bar{\psi}\chi Z_\chi\psi \; . \tag{3.28}$$

In order to compute χ for any specific fermion one, of course, must project out left-handed and right-handed components: $\psi=\psi_L+\psi_R$.

C. Third Z in E_6: Z_ψ

An example of a Z coupling with equal strength to all known light fermions may be obtained by decomposing

$$E_6\supset SO(10)\times U(1)_\psi \; . \tag{3.29}$$

The fundamental representation 27 of E_6, to which left-handed fermions are assumed to belong, breaks down according to this decomposition as

$$\underline{27}=\underline{16}\left[\frac{1}{\sqrt{24}}\right]+\underline{10}\left[\frac{-2}{\sqrt{24}}\right]+\underline{1}\left[\frac{4}{\sqrt{24}}\right] , \tag{3.30}$$

where the numbers in parentheses denote the $U(1)_\psi$ charges Q_ψ normalized the same way as other generators. For a 27-plet we have

$$\sum I_{3L}{}^2=\sum I_{3R}{}^2=3 \tag{3.31}$$

and so we stipulate $\sum Q_\psi{}^2=3$ as well.

We assume the known *light* fermions all belong to the SO(10) 16-plet. Since $Q_\psi(\underline{16})=1/\sqrt{24}$, all coupling strengths (to those left-handed fermions and antifermions) are equal.[34] This corresponds to purely axial-vector coupling. (The photon, which interacts vectorially, couples with equal and opposite strengths to left-handed fermions and antifermions.)

The fermions belonging to the SO(10) 10-plets and singlets in Eq. (3.30) are presumed heavier than those in the 16-plet. Their ψ charges are greater, and so the Z_ψ will decay more readily to them. In fact, the Z_ψ decays to ordinary fermions (including the right-handed neutrino) only $\frac{2}{9}$ of the time. The branching ratios of Z_ψ are summarized in Table I. (We have assumed that Z_ψ is much heavier than any of the fermions). If there are just three generations of fermions, $B(Z_\psi\to e^+e^-)=\frac{1}{108}\simeq 1\%$.

TABLE I. Branching ratios per generation of a hypothetical Z boson coupled to the charge Q_ψ of U(1) in $E_6\supset SO(10)\times U(1)_\psi$. Here Q is a charge $-\frac{1}{3}$ heavy quark, and L are heavy leptons.

SO(10):	16-plet				10-plet			Singlet
$u\bar{u}$	$d\bar{d}$	$e\bar{e}$	$\nu\bar{\nu}$	$N\bar{N}$	$Q\bar{Q}$	L^+L^-	L^0L^0	$L^0\bar{L}^0$
$\frac{1}{12}$	$\frac{1}{12}$	$\frac{1}{36}$	$\frac{1}{72}$	$\frac{1}{72}$	$\frac{1}{3}$	$\frac{1}{9}$	$\frac{1}{9}$	$\frac{2}{9}$
		$\frac{2}{9}$				$\frac{5}{9}$		$\frac{2}{9}$

The Lagrangian for the coupling of Z_ψ to fermions is

$$\mathscr{L}_{Z_\psi}=-g_\psi\bar{\psi}Q_\psi Z_\psi\psi \; . \tag{3.32}$$

To estimate the coupling g_ψ, we shall assume the breakdown (3.29) occurs at the common unification point described earlier, so that all U(1) couplings are of the same strength: $g_\psi=g_\chi=g_Y$ [see Eq. (3.21)].

D. Neutral bosons with Yang-Mills couplings

(1) An example of neutral bosons with SU(2) couplings occur in a class of E_6 models considered in Ref. 33. These are the ones (class III of Ref. 33) in which

$$E_6\supset SU(2)_{\text{inert}}\times SU(6) \; . \tag{3.33}$$

The "inert" SU(2) factor does not contribute to the electric charge. The 27-plet breaks down into

$$\underline{27}=(2_I,6^*)+(1_I,15) \; , \tag{3.34}$$

where[33]

$$(2_I,6^*)=\begin{bmatrix} \bar{h}_1 & \bar{d}_1 \\ \bar{h}_2 & \bar{d}_2 \\ \bar{h}_3 & \bar{d}_3 \\ \nu_E & \nu \\ E^- & e^- \\ \bar{N} & n^0 \end{bmatrix}_L , \tag{3.35}$$

$$(1_I,15)=\begin{bmatrix} 0 & \bar{u}_3 & \bar{u}_2 & d_1 & u_1 & h_1 \\ & 0 & \bar{u}_1 & d_2 & u_2 & h_2 \\ & & 0 & d_3 & u_3 & h_3 \\ & & & 0 & e^+ & \bar{N}_E \\ & & & & 0 & E^+ \\ & & & & & 0 \end{bmatrix}_L . \tag{3.36}$$

There is a triplet of neutral gauge bosons associated with $SU(2)_I$ (I for inert). These couple only to the doublet members in Eq. (3.35). The $I_{3I}=0$ boson we shall call Z_I^3, and the $I_{3I}=\pm 1$ bosons may be denoted

$$Z_I^{(\pm)} \equiv (Z_I^1 \mp i Z_I^2)/\sqrt{2} \; ,$$

recalling that they are electromagnetically neutral.

If Z_I^3 is much heavier than any of the fermions, its branching ratios to $d\bar{d}$ and e^+e^- are $\frac{1}{4}$ and $\frac{1}{12}$ per generation (the remaining decays would be to neutrinos and exotic fermions), or $\frac{1}{12}$ and $\frac{1}{36}$ for three generations.

The coupling of Z_I^3 to fermions may be obtained from the Lagrangian

$$\mathscr{L}_{Z_I} = -g_I \bar{\psi} \frac{\vec{\tau}_I}{2} \cdot \vec{Z}_I \psi \; , \tag{3.37}$$

where ψ denotes the fermions in (3.35). The coupling g_I evolves according to the behavior of an SU(2) from the point at which the breakdown (3.33) occurs. We assume for simplicity that this point is the same at which weak SU(2) splits off from other subgroups of the unifying group, and thus shall take

$$g_I{}^2 = g_2{}^2 = 8 G_F M_W{}^2/\sqrt{2} \; . \tag{3.38}$$

The $Z_I^{(\pm)}$ connect ordinary fermions with exotic ones, and thus cannot be produced via ordinary Drell-Yan processes. However, one could imagine situations in which the boson Z_I^3 is heavy enough to decay to one real and one virtual, or two real, $Z_I^{(\pm)}$. The branching ratio mentioned above then would be lowered. Correspondingly, however, exotic fermions might play a more significant role in Z_I^3 final states.[35]

(2) Other neutral bosons with Yang-Mills couplings have been proposed in connection with the structure of fermion generations.[36] Normally one expects masses of "horizontal" (generation-changing) bosons to be above 10—100 TeV.[37] It is possible to avoid these bounds for certain "diagonal" bosons Z_D, which then are subject only to the bound $M_{Z_D} \gtrsim 3 M_{Z_0}$ from neutral-current interactions. Such bosons may couple differently to different generations. In the simplest models[36] they would couple vectorially.

As an example, let us consider a horizontal $\mathrm{SU(3)}_H$ symmetry, acting over e,μ,τ generations with the same coupling strength as $\mathrm{SU(3)}_{\mathrm{color}}$, which we take to be $\alpha_H \equiv g_H{}^2/4\pi = \alpha_S \simeq 0.1$ for present purposes. The Lagrangian is

$$\mathscr{L}_{Z_H} = -g_H \bar{\psi} \frac{\bar{\lambda}}{2} \cdot \vec{Z}_H \psi \; . \tag{3.39}$$

Let us assume that only Z_H^8 is light. Then its branching ratios to $e^+e^-, \mu^+\mu^-$, and $\tau^+\tau^-$ would be, respectively, $\frac{1}{48}$, $\frac{1}{48}$, and $\frac{1}{12}$. In order for $\mathrm{SU(3)}_H$ to be anomaly free, all couplings must be vectorlike, so that a right-handed neutrino N is required. The branching ratios of Z_H to $N_e\bar{N}_e$, $N_\mu \bar{N}_\mu$, and $N_\tau \bar{N}_\tau$ are predicted to be $\frac{1}{96}$, $\frac{1}{96}$, and $\frac{1}{24}$.

IV. PRODUCTION CROSS SECTIONS

A Drell-Yan process for production of a vector boson for mass M by colliding hadrons A and B has the cross section per unit rapidity[38]

$$\frac{d\sigma}{dy} = \frac{4\pi^2 x_1 x_2}{3M^3} \sum_{i,j} f_i^{(A)}(x_1) f_j^{(B)}(x_2) \Gamma_{ij} \; , \tag{4.1}$$

where $f_i^{(A,B)}$ are the structure functions of quark i in hadrons A and B, and Γ_{ij} is the partial width of the vector boson into the quark pair $q_i q_j$. The momentum fractions x_1 and x_2 are related to the rapidity y by

$$\left.\begin{matrix} x_1 \\ x_2 \end{matrix}\right\} = (M/\sqrt{s}\,) e^{\pm y} \; , \tag{4.2}$$

where $\sqrt{s}$ is the total c.m. energy.

With the couplings determined in Sec. III (the Z_0 partial widths are also quoted for convenience), we calculate ($x_W \equiv \sin^2\theta$)

$$\Gamma\left[W_{L,R}^{\pm} \to \begin{Bmatrix} u\ \bar{d} \\ d\ \bar{u} \end{Bmatrix} \right] = \frac{G_F M_W{}^2}{\sqrt{2}} \frac{M_{W_{L,R}}}{2\pi} \cos^2\theta_C \tag{4.3}$$

(here θ_C is the Cabibbo angle, $\cos^2\theta_C = 0.948$[39]);

$$\Gamma(Z_0 \to u\bar{u}) = \frac{G_F M_{Z_0}{}^2}{\sqrt{2}} \frac{M_{Z_0}}{4\pi} [(1-\tfrac{4}{3}x_W)^2 + (\tfrac{4}{3}x_W)^2] \; , \tag{4.4}$$

$$\Gamma(Z_0 \to d\bar{d}) = \frac{G_F M_{Z_0}{}^2}{\sqrt{2}} \frac{M_{Z_0}}{4\pi} [(-1+\tfrac{2}{3}x_W)^2 + (\tfrac{2}{3}x_W)^2] \; , \tag{4.5}$$

$$\Gamma(Z_\chi \to u\bar{u}) = \tfrac{1}{5}\Gamma(Z_\chi \to d\bar{d}) = \frac{G_F M_{Z_0}{}^2}{\sqrt{2}} \frac{M_{Z_\chi}}{4\pi} \frac{x_W}{3} \; , \tag{4.6}$$

$$\Gamma(Z_\psi \to u\bar{u}) = \Gamma(Z_\psi \to d\bar{d}) = \frac{G_F M_{Z_0}{}^2}{\sqrt{2}} \frac{M_{Z_\psi}}{4\pi} \frac{5}{9} x_W \; , \tag{4.7}$$

$$\Gamma(Z_I^3 \to u\bar{u}) = 0 \; , \tag{4.8}$$

$$\Gamma(Z_I^3 \to d\bar{d}) = \frac{G_F M_W{}^2}{\sqrt{2}} \frac{M_{Z_I^3}}{4\pi} \; , \tag{4.9}$$

$$\Gamma(Z_H^8 \to u\bar{u}) = \Gamma(Z_H^8 \to d\bar{d}) = \frac{\alpha_H}{12} M_{Z_H^8} \; . \tag{4.10}$$

It is amusing to compare the partial widths for Z_0 and Z_χ decays. For $\sin^2\theta = 0.22$, we find

$$\frac{\Gamma(Z_\chi \to u\bar{u})}{\Gamma(Z_0 \to u\bar{u})} \simeq \frac{1}{8} \frac{M_{Z_\chi}}{M_{Z_0}} \; , \tag{4.11}$$

$$\frac{\Gamma(Z_\chi \to d\bar{d})}{\Gamma(Z_0 \to d\bar{d})} \simeq \frac{1}{2} \frac{M_{Z_\chi}}{M_{Z_0}} \; . \tag{4.12}$$

The Z_χ is hard to produce in hadron-hadron collisions because of these small couplings.

The branching ratios calculated in Sec. III may now be used to obtain values of $B\,d\sigma/dy$ in terms of quark distribution functions. The dimensionless quantities $sB\,d\sigma/dy$ are summarized in Table II. Thus, to illustrate,

TABLE II. Values of $sB\,d\sigma/dy$ for production and decay of specific gauge bosons. Here $U_A\bar{D}_B = U_A(X_1)\bar{D}_B(X_2)$, etc.

Boson	Final state	B (%)	Value of $sB\,d\sigma/dy$ Coefficient of $U_A\bar{D}_B+\bar{D}_AU_B(W^+)$ or $D_A\bar{U}_B+\bar{U}_AD_B(W^-)$	
$W^{\pm}$	$l\nu$	8.3	8.87×10^{-3}	
$W_R^{\pm}$	lN	8.3		
			Coefficient of $U_A\bar{U}_B+\bar{U}_AU_B$	Coefficient of $D_A\bar{D}_B+\bar{D}_AD_B$
Z_0		3.1	1.31×10^{-3}	1.67×10^{-3}
Z_χ		4.16	2.22×10^{-4}	1.11×10^{-3}
Z_ψ	l^-l^+	0.93	8.23×10^{-5}	8.23×10^{-5}
Z_I^3		2.78	0	1.57×10^{-3}
Z_H^8		2.08	2.28×10^{-3}	2.28×10^{-3}

$$B\frac{d\sigma}{dy}(A+B\to \underbrace{W^+}_{\to l^+\nu}+\cdots)=\frac{8.87\times10^{-3}}{s}(U_A\bar{D}_B+\bar{D}_AU_B)\ , \qquad (4.13)$$

where, for example, $U_A\bar{D}_B \equiv U_A(x_1)\bar{D}_B(x_2)$.

The quark distribution functions used to evaluate the cross sections in Table II are those of Ref. 6, based on CERN-Dortmund-Heidelberg-Saclay (CDHS) neutrino data, numerically extrapolated to large Q^2 using the Altarelli-Parisi equations with $\Lambda=0.29$ GeV.

The predicted values of $B\sigma$ for $W^{\pm}$ and Z_0 production are summarized in Table III. At $\sqrt{s}=0.54$ GeV, they are consistent with rates observed at the CERN $p\bar{p}$ collider.[3] We expect them to be 4 or 5 times larger at $\sqrt{s}=2$ TeV.

At $\sqrt{s}=2$ TeV, many other possibilities already exist for new-gauge-boson production. A sample of these is given in Fig. 1. These illustrate the range of production cross sections one can expect with gauge couplings. Cross sections for Z_I^3 and Z_H^8 production are not shown. They are bracketed by the range of those in Fig. 1. At a level of $B\sigma \geq 10^{-36}$ cm^2 (a reasonable lower limit for an experiment at the Fermilab Tevatron) the highest accessible masses range from $\lesssim 200$ GeV/c^2 for the Z_ψ to nearly 500 GeV/c^2 for either sign of W_R (and above 500 GeV/c^2 if both signs are combined).

For higher gauge-boson masses (above 200—500 GeV/c^2, depending on their couplings) one must employ c.m. energies above 2 TeV. For illustration we shall consider pp and $p\bar{p}$ collisions in the range $10 \leq \sqrt{s} \leq 40$ GeV. The cross sections times leptonic branching ratios ($l\nu$ or lN for W_R, l^-l^+ for Z's) for various gauge bosons are shown in Figs. 2—5.

An average luminosity of up to $\mathscr{L}=10^{33}$ cm^{-2} sec^{-1} is often quoted for planned large pp colliders. For $p\bar{p}$ colliders it may be possible to attain $\mathscr{L}=10^{32}$ cm^{-2} sec^{-1}. If one wishes to collect at least 10 events in an experiment of 10^7 sec, the lowest measurable values of σB are about 10^{-6} and 10^{-5} nb, respectively.

In Table IV we show maximum values of W_R and Z_χ masses attainable at various cross-section levels. A substantial jump above the 200—500-GeV/c^2 range attainable at the Tevatron is possible for $\mathscr{L}_{pp}=10^{33}$ cm^{-2} sec^{-1} or $\mathscr{L}_{p\bar{p}}=10^{32}$ cm^{-2} sec^{-1} and for $\sqrt{s}=40$ TeV. As an example, it may be possible to produce W_R up to about 9 TeV if values of $\sigma B=10^{-39}$ cm^2 can be measured in pp collisions. The corresponding value for $p\bar{p}$ collisions with

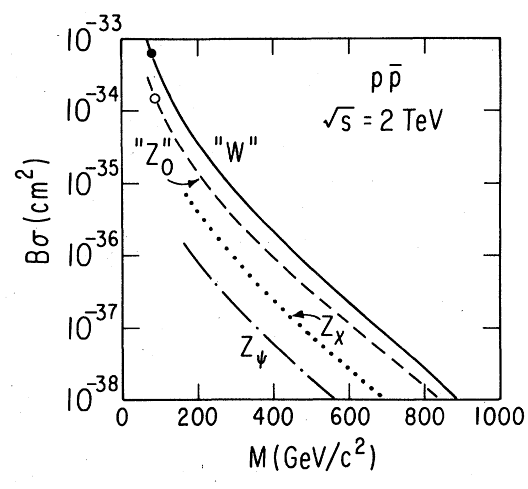


FIG. 1. Values of leptonic branching ratio B times cross section σ for production of various new heavy gauge bosons in $p\bar{p}$ collisions at $\sqrt{s}=2$ TeV, as functions of gauge-boson mass M. Solid line: W_R^+ or W_R^-. Dark circle: prediction for standard W. Dashed line: "Z_0" (couplings that of standard Z_0 but mass a free parameter). Open circle: prediction for standard Z_0. Dotted line: Z_χ [neutral boson in SO(10)/SU(5)]. Dash-dotted line: Z_ψ [neutral boson in E$_6$/SO(10)].

TABLE III. Values of $B\sigma$ in nb for production of $W^{\pm}$ (81 GeV/c^2) and Z_0 (92 GeV/c^2) in $p\bar{p}$ reactions. Here $B_W=8.3\%$, $B_Z=3.1\%$.

$\sqrt{s}$ (TeV)	$W^{\pm}$	Z_0
0.54	0.127	0.031
2.0	0.64	0.15

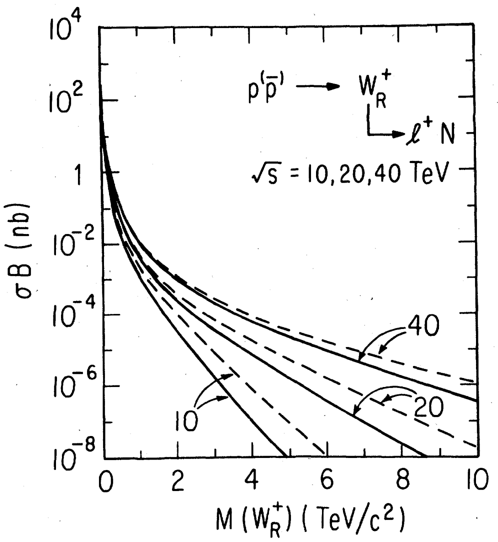


FIG. 2. Cross sections σ times branching ratio B to eN (8.3%) for W_R^+ in pp (solid curves) and $\bar{p}p$ (dashed curves) collisions. The three sets of curves refer to $\sqrt{s}=10$, 20, and 40 TeV, in ascending order.

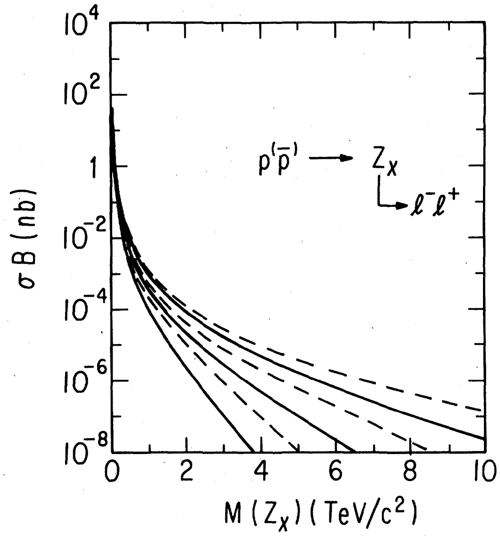


FIG. 4. Same as Fig. 2 for Z_X, with decay branching ratio replaced by $B_{ee}=4.16\%$.

$\sigma B=10^{-38}$ cm^2 is about 7 TeV.

We now display some rapidity distributions for light and heavy gauge bosons.

In Fig. 6 we compare rapidity distributions for $p\bar{p}\rightarrow W^+ + \cdots \rightarrow l^+ + \cdots$ at $\sqrt{s}=0.54$, 2, 10, and 40 TeV. One can see that the cross section builds up to a large extent at high y values, where W bosons are difficult to detect, but that an appreciable increase of production still occurs in the central region.

At $\sqrt{s}=40$ TeV, the effect of valence quarks can be clearly seen in W^+ production (Fig. 7). Two "wings" occur for pp collisions, and one wing for $p\bar{p}$. The wings are absent for $pp\rightarrow W^-$.

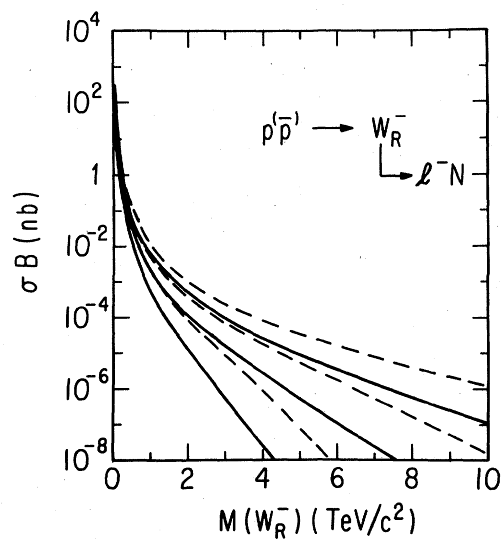


FIG. 3. Same as Fig. 2 for W_R^-.

A similar comparison for Z_0 production is made in Fig. 8. (By $\sqrt{s}=40$ TeV, there is almost no difference between total cross sections for pp and $p\bar{p}\rightarrow Z_0$). Substantial production is occurring at very high $|y|$ values.

Examples of heavy boson production at $\sqrt{s}=40$ TeV are given for W_R of mass 2 TeV/c^2 and Z_X of 1 TeV/c^2 in Figs. 9 and 10. The maximum rapidity for production of any boson of mass M is

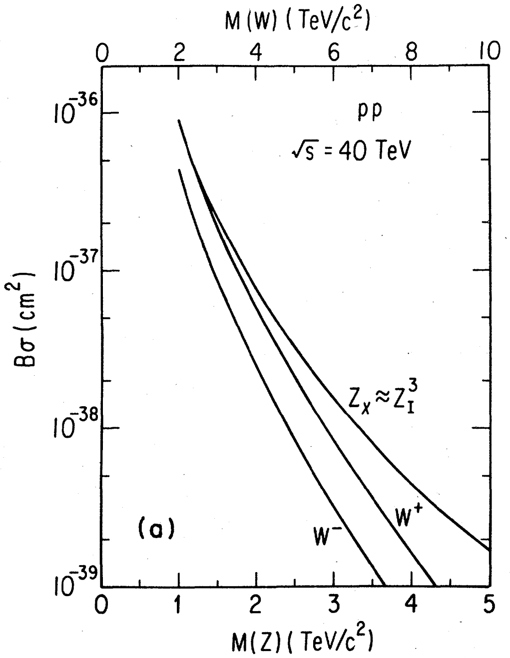
$$|y|_{\max}=\ln(\sqrt{s}/M) . \tag{4.14}$$

For masses above 1 TeV, the detection of heavy gauge bosons even at $\sqrt{s}=40$ TeV does not require the observation of especially high rapidities, since $\ln 40=3.7$. Moreover, events are fairly uniformly distributed in y.

V. FORWARD-BACKWARD ASYMMETRIES

A distinctive signature of the W, Z, and other hypothetical heavier gauge bosons is the forward-backward asymmetries in their production and leptonic decays. Consider the two-body process $q\bar{q}\rightarrow l\bar{l}$ in the quark-antiquark center of mass. Let the angle between l and q be θ^* in this frame. Then the angular distribution is a linear combination of $(1+\cos\theta^*)^2$ and $(1-\cos\theta^*)^2$ contributions:

$$\frac{d\sigma}{d(\cos\theta^*)} \propto (L_q{}^2L_l{}^2+R_q{}^2R_l{}^2)(1+\cos\theta^*)^2 + (L_q{}^2R_l{}^2+R_q{}^2L_l{}^2)(1-\cos\theta^*)^2 , \tag{5.1}$$



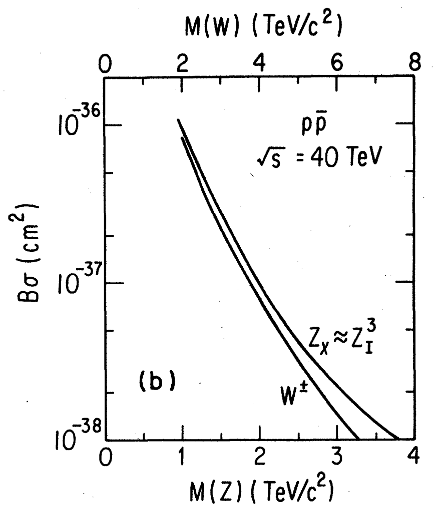


FIG. 5. $B\sigma$ for heavy W and Z production at $\sqrt{s}=40$ TeV. Branching ratios are as in Figs. 2–4, and $B_{ee}(Z_I^3)=2.78\%$. (a) pp collisions; (b) $p\bar{p}$ collisions.

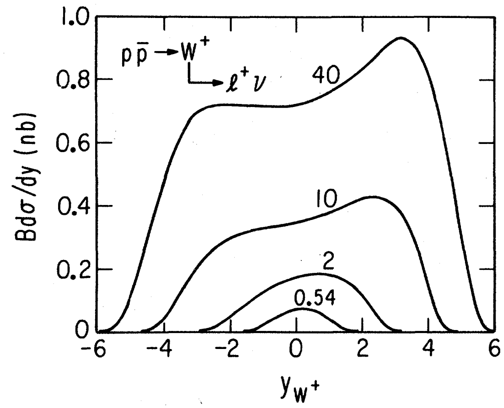


FIG. 6. Rapidity distributions for $p\bar{p}\to W^+ + \cdots \to l^+ + \cdots$ at $\sqrt{s}=0.54$, 2, 10, and 40 TeV (labels on curves).

where L_q, L_l are left-handed couplings to quarks and leptons, and R_q, R_l are right-handed couplings. For $A+B\to W_{L,R}^+ + \cdots \to l^+ + \cdots$,

$$\frac{d^2\sigma}{dy\,d(\cos\theta^*_{l^+})} \propto U_A(x_1)\bar{D}_B(x_2)(1-\cos\theta^*)^2 + \bar{D}_A(x_1)U_B(x_2)(1+\cos\theta^*)^2 \; ; \quad (5.2)$$

for $A+B\to W_{L,R}^- + \cdots \to l^- + \cdots$,

$$\frac{d^2\sigma}{dy\,d(\cos\theta^*_{l^-})} \propto D_A(x_1)\bar{U}_B(x_2)(1+\cos\theta^*)^2 + \bar{U}_A(x_1)D_B(x_2)(1-\cos\theta^*)^2 \; . \quad (5.3)$$

For Z production, $A+B\to Z+\cdots \to l^- + \cdots$,

TABLE IV. Maximum $W_R^\pm$ and Z_χ masses in TeV/c^2 attainable at specific cross-section levels.

	$\sqrt{s}$ (TeV)	pp $B\sigma=10^{-39}$ cm^2	10^{-38} cm^2	$\bar{p}p$ 10^{-38} cm^2	10^{-37} cm^2
W_R^+	10	3.2	2.3	2.9	1.9
	20	5.3	3.8	4.4	2.7
	40	8.6	5.8	6.5	3.7
W_R^-	10	2.7	2.0	same as W_R^+	
	20	4.5	3.1		
	40	7.3	4.8		
Z_χ	10	2.2	1.5	1.9	1.1
	20	3.5	2.3	2.8	1.5
	40	5.6	3.3	3.8	2.0

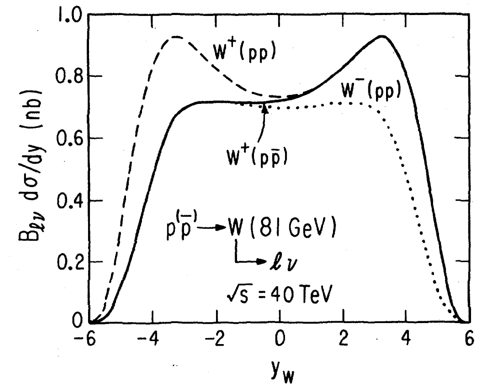


FIG. 7. Rapidity distributions for W^+ production by pp and $p\bar{p}$, and for W^- by pp, at $\sqrt{s}=40$ TeV. Note that $B[d\sigma(y)/dy]\,(p\bar{p}\to W^-)=B[d\sigma(-y)/dy](p\bar{p}\to W^+)$.

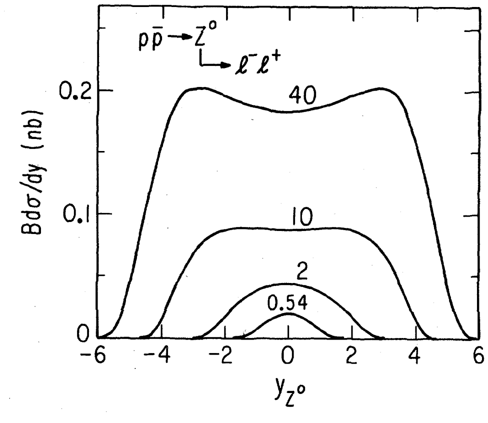


FIG. 8. Rapidity distributions for $p\bar{p}\to Z_0\to l^-l^+$ at $\sqrt{s}=0.54$, 2, 10, and 40 TeV (labels on curves).

$$\frac{d^2\sigma}{dy\,d(\cos\theta^*_{l^-})}\propto U_A(x_1)\bar{U}_B(x_2)[(L_u{}^2L_l{}^2+R_u{}^2R_l{}^2)(1+\cos\theta^*)^2+(L_u{}^2R_l{}^2+R_u{}^2L_l{}^2)(1-\cos\theta^*)^2]$$

$$+\bar{U}_A(x_1)U_B(x_2)[(L_u{}^2L_l{}^2+R_u{}^2R_l{}^2)(1-\cos\theta^*)^2+(L_u{}^2R_l{}^2+R_u{}^2L_l{}^2)(1+\cos\theta^*)^2]+(U\to D)\ , \qquad (5.4)$$

where in (5.2)—(5.4) we denote the forward direction by that in which hadron A is traveling. We shall always take A to be a proton; B may be p or $\bar{p}$.

Forward-backward asymmetries $A_{\rm FB}$ may be constructed by binning data with respect to rapidity. For any fixed y, we define

$$A_{\rm FB}=\frac{F-B}{F+B}\ , \qquad (5.5)$$

where

$$F\pm B=\left[\int_0^1\pm\int_{-1}^0\right]d(\cos\theta^*)\frac{d^2\sigma}{dy\,d(\cos\theta^*)}\ . \qquad (5.6)$$

We then find, for $A+B\to W^+_{L,R}+\cdots\to l^++\cdots$,

$$A_{\rm FB}=\tfrac{3}{4}\frac{\bar{D}_AU_B-U_A\bar{D}_B}{\bar{D}_AU_B+U_A\bar{D}_B}\ ; \qquad (5.7)$$

for $A+B\to W^-_{L,R}+\cdots\to l^-+\cdots$,

$$A_{\rm FB}=\tfrac{3}{4}\frac{D_A\bar{U}_B-\bar{U}_AD_B}{D_A\bar{U}_B+\bar{U}_AD_B}\ ; \qquad (5.8)$$

and for $A+B\to Z+\cdots\to l^-+\cdots$,

$$F\pm B\propto\begin{Bmatrix}\tfrac{4}{3}\\1\end{Bmatrix}[(U_A\bar{U}_B\pm\bar{U}_AU_B)(L_u{}^2\pm R_u{}^2)(L_l{}^2\pm R_l{}^2)$$

$$+(U\to D)]\ . \qquad (5.9)$$

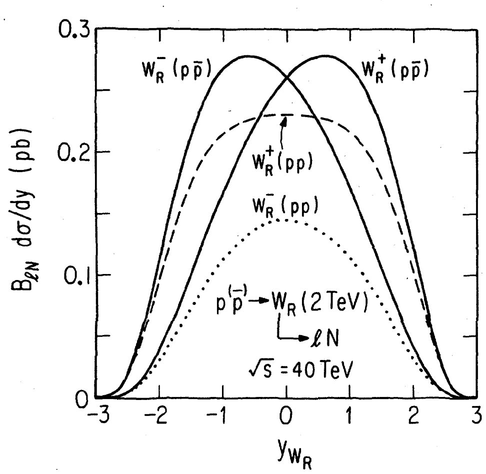


FIG. 9. Rapidity distributions for $W_R^{\pm}$ (2 TeV) production by pp (dashed and dotted curves) and $p\bar{p}$ (solid curves) at $\sqrt{s}=40$ TeV.

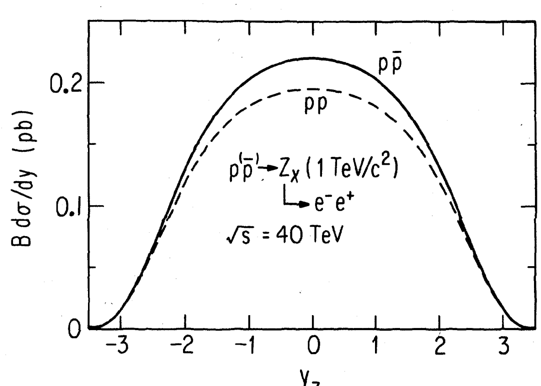


FIG. 10. Rapidity distributions for Z_χ (1 TeV) production by pp (dashed curve) and $p\bar{p}$ (solid curve) at $\sqrt{s}=40$ TeV.

TABLE V. Relative values of left-handed and right-handed couplings to quarks and leptons for various Z's (defined in Sec. III). Here we take $x_W=0.22$, and normalize $L_u^2+R_u^2=1$.

	u quarks		d quarks		charged leptons	
Boson	L_u^2	R_u^2	L_d^2	R_d^2	L_e^2	R_e^2
Z_0	0.853	0.147	0.971	0.029	0.618	0.382
Z_χ	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{10}$	$\frac{9}{10}$	$\frac{9}{10}$	$\frac{1}{10}$
Z_ψ	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$
Z_I^3			0	1	1	0
Z_H^8	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$

What makes $A_{\rm FB}$ so useful for identifying gauge bosons is that each boson has a particular signature. When valence quarks dominate, in $p\bar{p}\to W^+ + \cdots \to l^+ + \cdots$, $U_A\bar{D}_B \gg \bar{D}_A U_B$, so $A_{\rm FB}\approx -\frac{3}{4}$, while in $p\bar{p}\to W^- + \cdots \to l^- + \cdots$, $D_A\bar{U}_B \gg \bar{U}_A D_B$, so $A_{\rm FB}\approx +\frac{3}{4}$. Positive leptons follow the $\bar{p}$ direction, while negative leptons follow the p direction. This conclusion does not depend on the handedness of the W's, as long as it is the same for quark and lepton couplings.

For Z's the relative values of left-handed and right-handed couplings are summarized in Table V. The left-right asymmetries for Z_0 are expected to be small and very sensitive to x_W because $L_e^2=R_e^2$ when $x_W=\frac{1}{4}$, a value close to the observed one. Larger asymmetries are expected for Z_χ, still diluted somewhat because $L_u^2=R_u^2$. No asymmetries are expected for Z_ψ (purely axial couplings) or Z_H^8 (purely vector couplings). Maximal asymmetries are expected for Z_I^3, which couples only to right-handed d and left-handed e^-.

We give examples of predicted asymmetries as functions of y in Figs. 11–16. Several comments are in order.

(1) In pp interactions, the forward-backward asymmetries are necessarily antisymmetric about $y=0$. They rise (or fall) to very near the maximum (or minimum) values attained in $p\bar{p}$ collisions at the same energy when one moves away from $y=0$.

(2) For $M/\sqrt{s} \gtrsim 0.1$, asymmetries in $p\bar{p}$ collisions are nearly constant over a wide y range. As $\sqrt{s}$ increases, these asymmetries become washed out in the central region as sea-sea collisions begin to play a larger role in gauge-boson formation.

(3) The asymmetries in Z production indeed are very sensitive to the specific forms of couplings, as one sees by comparing Figs. 13, 15, and 16. The asymmetries in Z_χ production fall off for large $|y|$ because the Z_χ is mainly produced by u quarks in this region, and $L_u^2=R_u^2$. The signs of the expected asymmetries in Z_0 production (for $x_W<\frac{1}{4}$) are opposite to those in Z_χ or Z_I^3 production, for the same beam and value of y.

If data can be binned with respect to rapidity, there should be no problem in determining asymmetries in $p\bar{p}$ or pp collisions. This method has been examined in detail for a specific detector design,[40] and found satisfactory for measuring asymmetries in the decay of a Z_χ with mass 1 TeV/c^2.

There are more global variables sensitive to weak asymmetries in W and Z decays. The detection of asymmetries in W decays is straightforward in $p\bar{p}$ interactions. As mentioned, positive leptons tend to follow the $\bar{p}$ direction, while negative leptons follow the p direction. This effect has already been used to measure the spin of the W.[41]

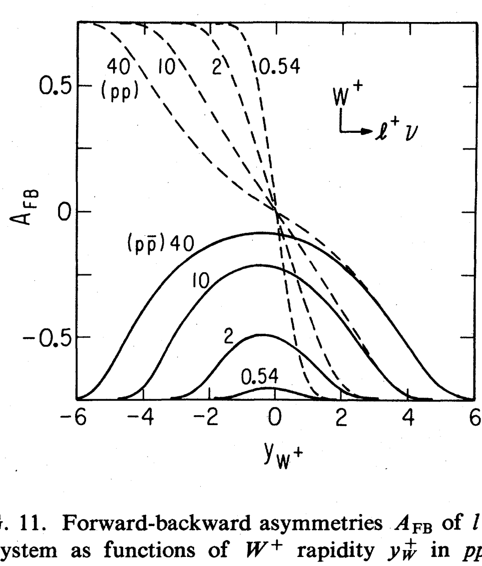


FIG. 11. Forward-backward asymmetries $A_{\rm FB}$ of l^+ in W^+ c.m. system as functions of W^+ rapidity y_W^+ in pp (dashed curves) and $p\bar{p}$ (solid curves) $\to W^+\to l^+\nu$. Curves are labeled by total c.m. energy in TeV.

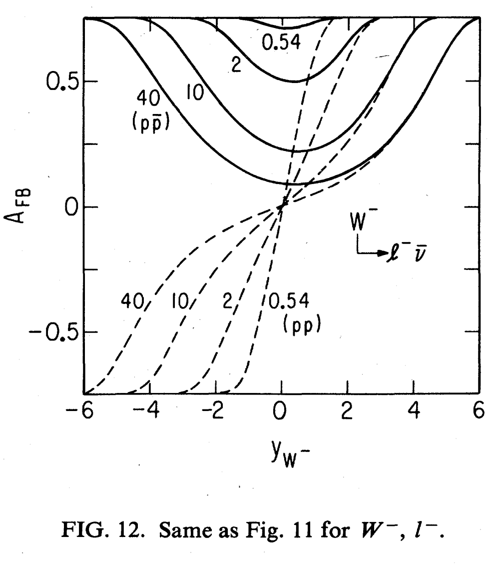


FIG. 12. Same as Fig. 11 for W^-, l^-.

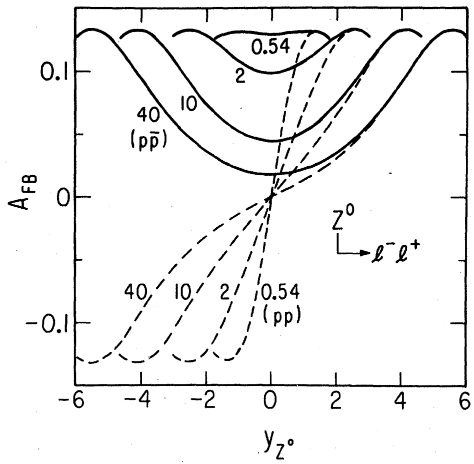


FIG. 13. A_{FB} of l^- vs y_{Z_0} in pp (dashed curves) and $p\bar{p}$ (solid curves)$\to Z_0 \to l^- l^+$. Curves are labeled by total c.m. energy in TeV.

The asymmetry is substantial, as shown in Fig. 17. It is expected to persist as long as valence quarks continue to play a major role in W production. As an example, we show the corresponding distribution for $p\bar{p}\to W_R^+(2$ TeV$)\to l^+$ at $\sqrt{s}=40$ TeV in Fig. 18. Here we have defined

$$\hat{x}^{\perp} \equiv p^{\perp}/p^{\perp}_{\max} = 2p^{\perp}/M = \sin\theta^* \ , \tag{5.10}$$

where M is the mass of the gauge boson. (As elsewhere, transverse momentum in gauge-boson production is neglected.)

In pp interactions, the lepton distributions are symmetric about 90°. However, the forward-backward asymmetries in $\cos\theta^*$ lead to different distributions in laboratory angles and transverse momenta for l^- and l^+. An example is shown in Fig. 19, which is a composite of two figures symmetric about 90°.

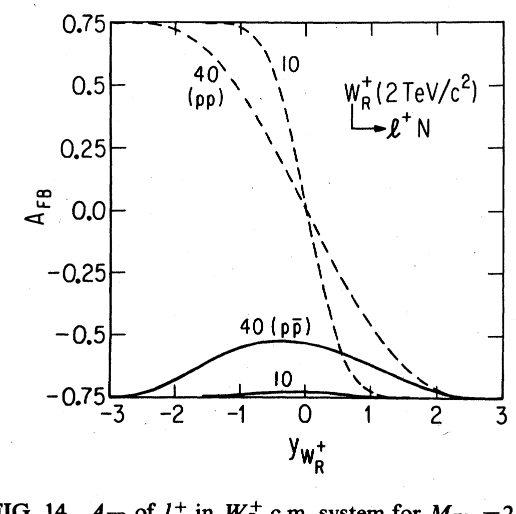


FIG. 14. A_{FB} of l^+ in W_R^+ c.m. system for $M_{W_R}=2$ TeV, vs $y_{W_R^+}$. (Other labels as in Fig. 11.) For W_R^-, reverse signs of A_{FB} and y_W.

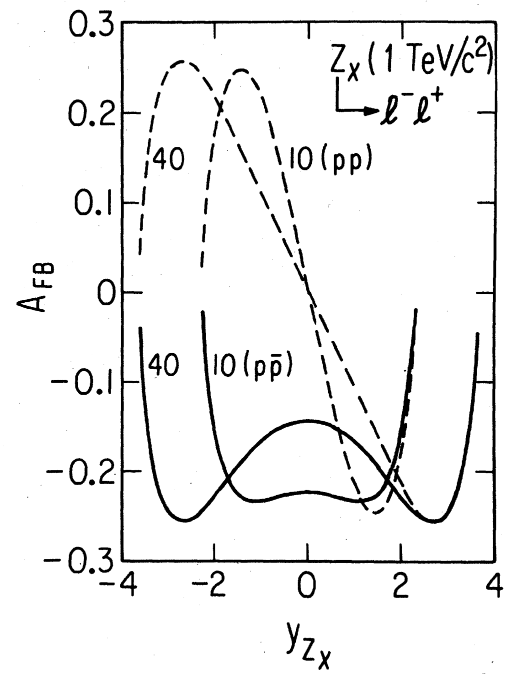


FIG. 15. A_{FB} of l^- vs y_{Z_χ} for $M_{Z_\chi}=1$ TeV$/c^2$, vs y_{Z_χ}. (Other labels as in Fig. 13.)

A global variable of some usefulness in pp interactions is the total lepton energy $\langle E_l \rangle$, where (here y is the boson rapidity)

$$E_l = \frac{M\hat{x}^{\perp}}{2\sin\theta} = \frac{M}{2}(\cosh y + \cos\theta^* \sinh y) \ . \tag{5.11}$$

For $pp \to W_{(R)}^+ \to l^+$, the lepton is thrown toward the cen-

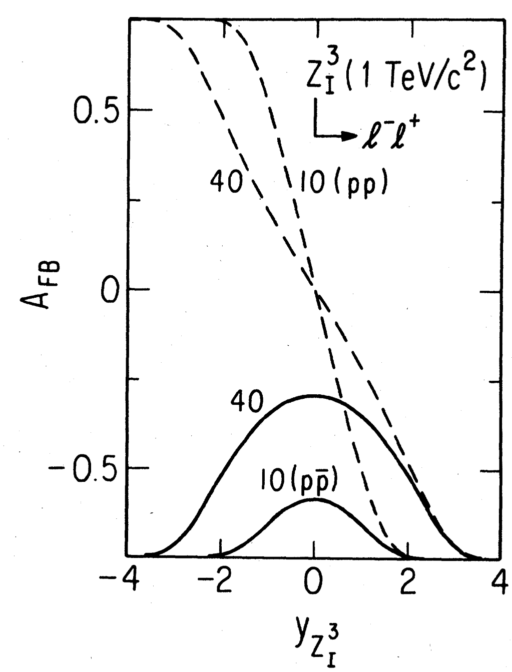


FIG. 16. Same as Fig. 15 for Z_I^3.

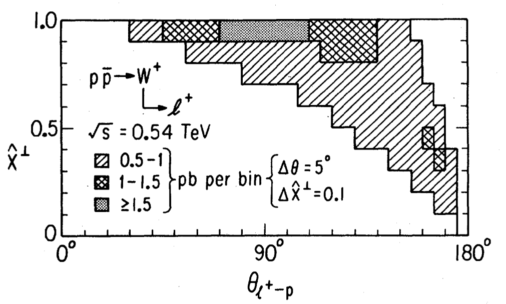


FIG. 17. Distribution of leptons l^+ in $p\bar{p} \to W^+ + \cdots \to l^+ + \cdots$ at $\sqrt{s} = 0.54$ TeV, with respect to laboratory angle $\theta_{l^+ - p}$ between l^+ and p, and the variable $\hat{x}^{\perp} = p^{\perp}(l^+)/p^{\perp}_{\max}(l^+)$. (1 pb $= 10^{-36}$ cm^2).

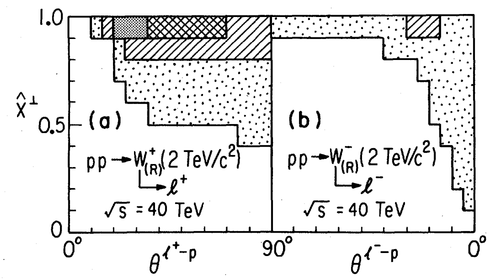


FIG. 19. Distribution of leptons l^+ (a) and l^- (b) in $pp \to W^{\pm}_R$ (2 TeV)$+ \cdots \to l^{\pm} + \cdots$ at $\sqrt{s} = 40$ TeV. Shading codes are as in Fig. 18.

tral region and its average energy is reduced. For $pp \to W^-_{(R)} \to l^-$, the lepton is thrown toward larger $|y|$ (smaller angles with respect to the beams), and has larger energy. For $pp \to W^{\pm}$, the lepton energies also differ by virtue of the differences in y distributions of W^+ and W^-, and are somewhat dependent on details of structure functions. However, for $pp \to Z$, the lepton energies are unequal purely because of the decay asymmetry.

Lepton distributions in Z decays are shown in Figs. 20—22. The patterns are not grossly asymmetric. Nonetheless, in $pp \to Z_{\chi} \to l^{\mp}$, the average lepton energies differ notably from one another.

Lepton energies are compared at $\sqrt{s} = 40$ TeV for various reactions in Fig. 23, as a function of gauge-boson mass. For $pp \to W^{\pm}_{(R)}$ (2 TeV), the $l^{\pm}$ energies differ by more than a factor of 2. By reversing the form of the coupling at the quark or lepton vertex ($V \pm A \to V \mp A$), we have changed that most of this asymmetry is indeed due to the coupling form, and not to the difference in y distributions of produced bosons.

The maximum mass of a particle for which the lepton energy asymmetry can be a useful variable depends on detector resolution as well as on the (calculable) distribution of lepton energies. A crude estimate in which these distributions are assumed to be very broad suggests that this method may be satisfactory even for Z_{χ} as high as 3 TeV/c^2 in mass at $\sqrt{s} = 40$ TeV, if a pp experiment with $\int \mathscr{L}\, dt = 10^{40}$ cm^{-2} can be performed. According to Fig. 5(a), one would expect 160 $Z_{\chi} \to l^- l^+$ decays under these conditions. The predicted energies are $\langle E_{l^-} \rangle = 1.94$ TeV, $\langle E_{l^+} \rangle = 2.49$ TeV for $M_{\chi} = 3$ TeV/c^2, $\sqrt{s} = 40$ TeV. This is a 25% difference. If 16 events are required to see such a difference at the 1σ level ($25\% = 1/\sqrt{16}$), 160 events will show more than a 3σ effect.

Another global variable sensitive to weak asymmetries is the forward-backward asymmetry of leptons in the gauge-boson center of mass, defined in terms of integrals over gauge-boson rapidity:

$$\langle A_{\mathrm{FB}} \rangle \equiv \int dy(F - B) \Big/ \int dy(F + B)\,, \tag{5.12}$$

with F and B defined in (5.6). This quantity is nonzero for $p\bar{p}$ interactions. Examples of its values are shown in Fig. 24.

As gauge-boson masses increase, asymmetries become more pronounced because valence quarks play a larger role in production. Cross sections decrease with mass, as shown in Fig. 5(b). The highest masses of gauge bosons for which an asymmetry can be detected as a 4σ effect are summarized in Table VI. A $p\bar{p}$ experiment at $\sqrt{s} = 40$ TeV, $\int \mathscr{L}\, dt = 10^{39}$ cm^2, can detect a 4σ asymmetry using this method up to $M_{Z_{\chi}} = 1.6$ TeV/c^2. The corresponding mass for pp, $\sqrt{s} = 40$ TeV, $\int \mathscr{L}\, dt = 10^{40}$ cm^2, using the energy asymmetry mentioned earlier, is 2.7 TeV/c^2 for a 4σ effect. However, as mentioned, that estimate was more dependent on detector details than the present one.

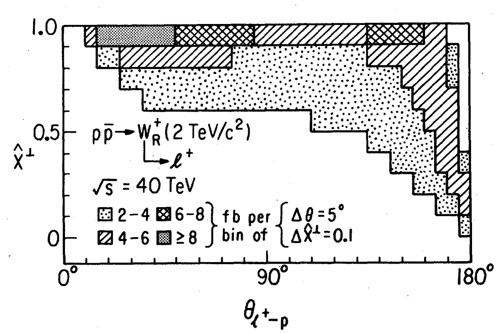


FIG. 18. Distribution of leptons l^+ in $p\bar{p} \to W^+_R$ (2 TeV) $+ \cdots \to l^+ + \cdots$ at $\sqrt{s} = 40$ TeV. (1 fb $= 10^{-39}$ cm^2).

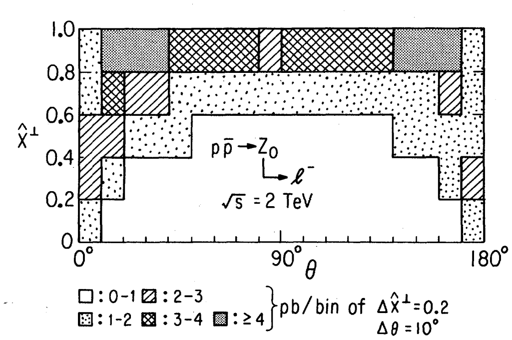


FIG. 20. Distribution of leptons l^- in $p\bar{p} \to Z_0 + \cdots \to l^- + \cdots$ at $\sqrt{s} = 2$ TeV.

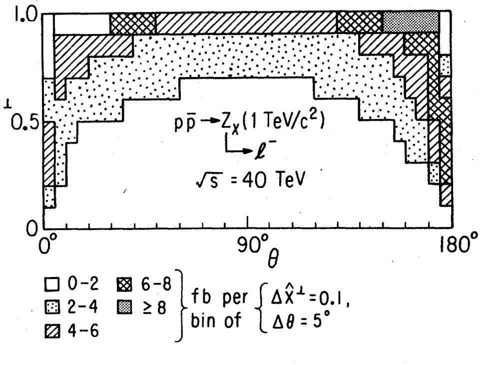


FIG. 21. Distribution of leptons l^- in $p\bar{p} \to Z_\chi$ (1 TeV)$+ \cdots \to l^- + \cdots$ at $\sqrt{s} = 40$ TeV.

For $p\bar{p} \to Z_0(92\ \text{GeV}/c^2) \to l^-$ at $\sqrt{s} = 2$ TeV, we find $\langle A_{FB} \rangle_y = 0.109$ at $x_W = 0.22$. The asymmetry is expected to vanish at $x_W = \frac{1}{4}$ and to be approximately linear in $x_W - \frac{1}{4}$ for small deviations from this value:

$$\langle A_{FB} \rangle_y \simeq 3.5(0.25 - x_W) . \tag{5.13}$$

The total cross section for $p\bar{p} \to Z_0 \to l^- l^+$ at $\sqrt{s} = 2$ TeV is estimated to be 150 pb (Table III). In an experiment with $\int \mathscr{L}\, dt = 10^{37}\ \text{cm}^{-2}$, we estimate that an asymmetry as small as $\langle A_{FB} \rangle_y = 0.07$ could be detected at the 4σ level. This would correspond to $x_W = 0.23$. The asymmetry is thus a potentially sensitive "vernier" for determining the deviation of x_W from $\frac{1}{4}$.

VI. EXPERIMENTAL SIGNATURES

A. Specific features of W_R and Z_χ decays

The detection of heavy gauge bosons is an experimental challenge that in our opinion has not been sufficiently addressed. A report in the 1982 Snowmass proceedings implies that the signatures for such bosons are straightforward.[42] For several reasons, we regard this optimism as premature.

(1) The leptons in W_R or Z decay can be quite energetic, with up to nearly 5 TeV of transverse energy and con-

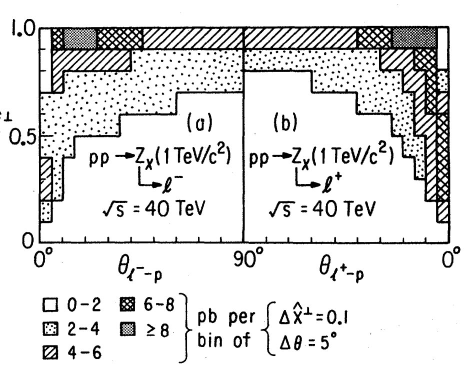


FIG. 22. Distribution of leptons l^- (a) and l^+ (b) in $pp \to Z_\chi$ (1 TeV)$+ \cdots \to l^- + \cdots$ at $\sqrt{s} = 40$ TeV. Shading codes as in Fig. 21.

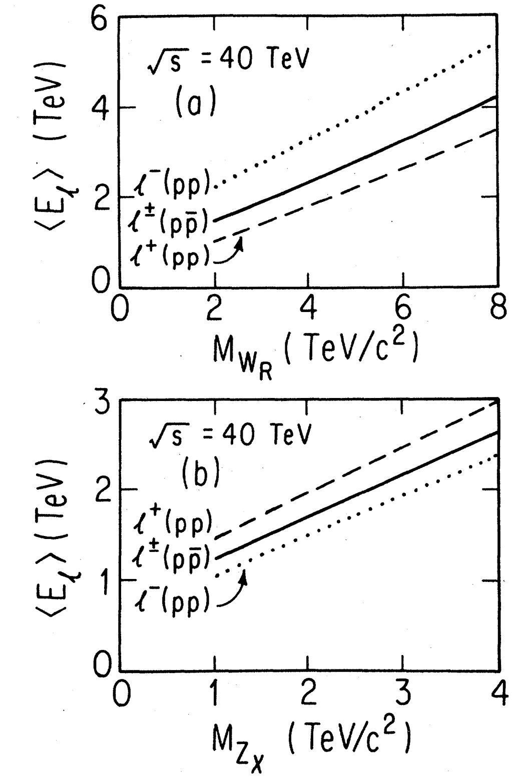


FIG. 23. Comparison of average lepton energies at $\sqrt{s} = 40$ TeV in $p\bar{p}$ collisions (solid curves) and pp collisions (l^+: dashed curves; l^-: dotted curves) for (a) W_R decay, (b) Z_χ decay.

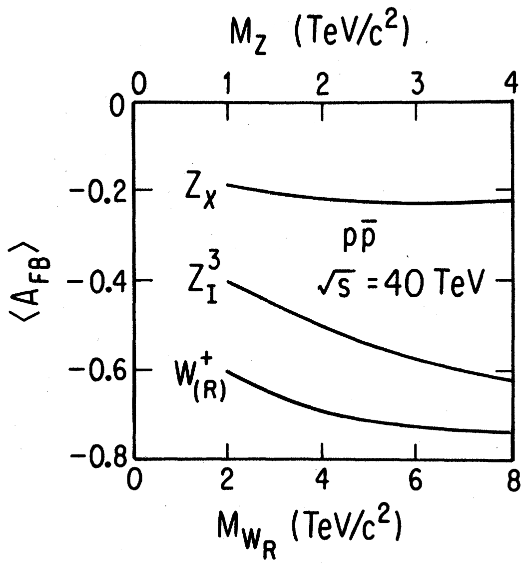


FIG. 24. Forward-backward asymmetry of leptons (l^- for Z production, l^+ for W^+ production) averaged over gauge-boson rapidity for $p\bar{p} \to Z_\chi$, Z_I^3, and W_R^+ at $\sqrt{s} = 40$ TeV. The forward direction is always defined by the proton.

TABLE VI. Examples of estimated highest masses for detection of a forward-backward asymmetry in gauge-boson decay. Bosons are produced in $p\bar{p}$ interactions at $\sqrt{s}=40$ TeV.

Gauge boson	Asymmetry	No. of events required	Maximum mass in TeV/c^2 for $\int \mathscr{L}\,dt=10^{39}$ cm^2
$W_R^{\pm}$	$\approx \mp 0.7$	20	5.6
Z_χ	≈ -0.2	200	1.6
Z_I^3	≈ -0.6	50	2.4

siderably more longitudinal energy (if $\sqrt{s}=40$ TeV). Muon identification at these energies is quite difficult. Electron identification may require very specific planning if the event rate is expected to be high, and particularly if more than one interaction per bunch crossing is anticipated.

(2) The leptonic decay of W_R involves the right-handed counterpart N of the neutrino. This neutral lepton may be considerably heavier than any of the quarks or leptons in its generation. Its mass is probably, though not necessarily, less than that of the W_R.[43] Its decay has distinctive signatures which we shall analyze presently.

(3) The decays of Z_χ to charged-lepton pairs are probably the most "ordinary" feature of this particle, since $B(Z_\chi \to e^+e^-)$ is predicted to be 4.2% for three fermion generations. Unusual features of Z_χ decay include the predicted relative suppression of $Q=\frac{2}{3}$ pairs [Eq. (4.6)] and the presence of a mode $Z_\chi \to N\bar{N}$ with branching ratio $\sim 10\%$ for each of three flavors of N (assuming three generations). As for W_R, the subsequent decay of N can provide a distinctive decay signature.

B. Some aspects of N decays

We now outline some properties of N that might be important in detecting W_R and Z_χ. More details can be found in Refs. 44 and 45.

(1) The N may be a Dirac particle, decaying to a "right-sign" lepton, or a Majorana particle capable of decaying to each sign of lepton.[44] *Lepton-sign information thus is crucial* in untangling signatures which might have originated in W_R or Z_χ decays.

(2) The N may decay via virtual W_R emission, via mixing with ν, or via both mechanisms.

If N decays via virtual W_R emission, its decay rate should be equal to

$$\Gamma(N_l \to l^- + \text{hadrons}) = 9\frac{G_F{}^2 M_N{}^5}{192\pi^3}\left[\frac{M_{W_L}}{M_{W_R}}\right]^4 , \quad (6.1)$$

where the factor of 9 refers to three colors times three flavors of weak isodoublet quarks $(u\bar{d}, c\bar{s}, t\bar{b})$. We thus expect

$$\tau_{N_l} = \frac{\tau_\mu}{9}\left[\frac{M_{W_R}}{M_{W_L}}\right]^4 \left[\frac{M_\mu}{M_N}\right]^5 . \quad (6.2)$$

This expression can be written as

$$\tau_{N_l} = (1.2\times 10^{-16}\ \text{sec})\left[\frac{M_{W_R}}{2\ \text{TeV}}\right]^4 \left[\frac{100\ \text{GeV}}{M_N}\right]^5 . \quad (6.3)$$

The leptonic decay modes, present for ordinary leptons such as τ, are absent here for the lightest N_l since the virtual W_R would have no lighter leptonic channel to which to decay. (A heavier N_l would also have leptonic decays, $N_{l'} \to l' + N_l + \bar{l}$.)

Even if decay occurs only via mixing with ν, the predicted decay rates grow with M_N. For $M_N \gg M_{W_L}$, we expect[45]

$$\Gamma(N_l \to l^- W^+) = \frac{G_F M_N{}^3}{8\pi\sqrt{2}} \, |U_{N_l}|^2 , \quad (6.4)$$

$$\Gamma(N_l \to \nu_l Z^0) = \frac{G_F M_N{}^3}{16\pi\sqrt{2}} \, |U_{N_l}|^2 . \quad (6.5)$$

The mixing parameter U_{N_l} is model-dependent. In a model in which N_l is a Majorana particle, we expect

$$U_{N_l} = m_D / M_{N_l} , \quad (6.6)$$

where m_D is a Dirac mass of order of m_l, the charged-lepton mass. Then

$$\Gamma_{\text{tot}}(N_l) = \frac{3G_F m_D{}^2 M_N}{16\pi\sqrt{2}} \quad (6.7)$$

or

$$\tau_{N_l} = (1.3\times 10^{-20}\ \text{sec})\left[\frac{1\ \text{GeV}}{m_D}\right]^2 \left[\frac{100\ \text{GeV}}{M_N}\right] . \quad (6.8)$$

It is conceivable that for m_D very small, even a N_l of mass 100 GeV can travel far enough to leave a detectable gap between its production and its decay.

If M_N is lighter than $M(W)$, Eq. (6.8) is replaced by[45,46]

$$\tau_{N_l} \simeq (5\times 10^{-12}\ \text{sec})\left[\frac{1\ \text{GeV}}{M_N}\right]^{5.2} |U_{N_l}|^{-2} . \quad (6.9)$$

The power of M_N in (6.9) reflects both the M^{-5} dependence in the Fermi expression for the lifetime, and an additional $M^{-0.2}$ dependence which is a phenomenological fit to the opening of new channels in virtual W and Z_0 decay. With U_{N_l} given by Eq. (6.6), we would then have

$$\tau_{N_l} \simeq (5\times 10^{-12}\ \text{sec})\left[\frac{1\ \text{GeV}}{M_N}\right]^{3.2} \left[\frac{1\ \text{GeV}}{m_D}\right]^2 . \quad (6.10)$$

Values of M_N above 1—2 GeV are not excluded by present experiments.[45] Thus, if N is sufficiently light, it could travel some distance from its production point before decaying, especially if produced in decays of a very heavy W_R or Z_χ.

It is thus crucial when searching for new heavy gauge bosons to anticipate the emission of neutral particles other than neutrinos which may have distinctive decay signatures. These neutral particles can (a) decay immediately,

(b) decay at a secondary vertex inside the detector, or (c) escape the detector entirely. The lifetime estimates (6.8)—(6.10) show that all three possibilities should be anticipated.

C. Rapidity distributions and asymmetries

We have discussed in some detail in Sec. V the distributions of leptons produced in heavy-gauge-boson decays. These can provide distinctive "fingerprints" of the bosons. Maximal asymmetries are expected for right-handed W's and for some types of Z's. The most likely new Z, the Z_χ, also is expected to have a characteristic asymmetry in its couplings to d quarks and electrons, though its coupling to u quarks does not exhibit this asymmetry.

For both pp and $p\bar{p}$ interactions, the asymmetry

$$A_{\rm FB}(y)=\frac{(d\sigma/dy)_{\cos\theta^*>0}-(d\sigma/dy)_{\cos\theta^*<0}}{(d\sigma/dy)_{\cos\theta^*>0}+(d\sigma/dy)_{\cos\theta^*<0}} \tag{6.11}$$

is a useful quantity. Here y is the rapidity of the gauge boson, while θ^* is the angle of lepton emission with respect to the proton direction in the gauge boson's rest frame. For pp interactions $A_{FB}(y)$ necessarily is antisymmetric about $y=0$.

Global variables which can be useful in detecting these asymmetries include the following.

(1) Total cross section in the hemisphere $\cos\theta^*>0$ vs its value for $\cos\theta^*<0$. This allows one to measure an asymmetry $\langle A_{\rm FB}\rangle$ which is nonzero for $p\bar{p}$ interactions.

(2) Lepton laboratory energies $\langle E_{l^-}\rangle$ vs $\langle E_{l^+}\rangle$. These will differ from one another for pp interactions. In the case of Z production, this difference is due entirely to the decay asymmetry. For $W^\pm$ production the difference between $\langle E_{l^-}\rangle$ and $\langle E_{l^+}\rangle$ contains information both on decay asymmetries and on differences in $d\sigma/dy$ for W^- and W^+.

There is some possibility for ambiguity in the sign of $\cos\theta^*$ if only a single lepton is observed. The two solutions correspond to different values of longitudinal momentum carried by the other lepton. If the other lepton is charged or decays rapidly (as N might), this ambiguity can be resolved directly. If it is not observed, its longitudinal momentum must be inferred by momentum balance in the detector.

D. Exotic products: general features

We have mentioned the Z_χ as one source of new particles: it decays frequently to pairs of new heavy leptons N, if $M_N<M_{Z_\chi}/2$. This is a general feature of new heavy gauge bosons. They can connect "old" states of matter with "new." A more dramatic example occurs for Z_ψ, which can decay to a whole new class of fermions [belonging to SO(10) 10-plet and singlets] as well as the more familiar SO(10) 16-plets. Possible signatures of these new states of matter include longer lifetimes (since they would mix weakly with the old states), characteristic weak-decay chains favoring lepton emission, and (of course) elevated masses (since the charged varieties have not been observed in e^+e^- annihilations).

E. Experimental design considerations

We conclude this section with some recommendation for experimental designs.

(1) It is crucial to determine lepton signs in searching for new gauge bosons. These are important for several reasons. First, one wishes to reduce backgrounds by comparing $l^\pm l^\pm$ pairs with l^+l^- pairs. Second, many of the asymmetry tests we suggest here depend crucially on lepton signs. In view of the large energy of the expected leptons, determination of their signs will be experimentally challenging at multi-TeV energies.

(2) Longitudinal-momentum balance often plays a role in resolving a two-fold kinematic ambiguity when a single lepton from a heavy gauge boson is detected. In order to resolve this ambiguity, nearly full coverage of 4π by the detector is desirable.

(3) We have not addressed the question of whether very heavy gauge bosons can be detected via their decays to quark-jet pairs. This appears to be a detailed question of detector resolution, but preliminary efforts to detect W and Z decays by this method so far have proved unpromising.

VII. SUMMARY AND DISCUSSION

We have given an overview of right-handed W ("W_R") and heavy Z("Z_χ", . . .) production and observation in multi-TeV pp and $\bar{p}p$ collisions. The crucial features are these:

(1) W_R should be produced very similarly to W_L, aside from its mass. Its leptonic decay is characterized by coupling to eN rather than $e\nu$, where N is the "right-handed neutrino."

(2) Z_χ production in pp and $\bar{p}p$ collisions is somewhat more difficult than that of a heavy Z_0 with standard couplings. (The Z_χ is the unique additional neutral boson expected if one gauges $B-L$ as well as weak hypercharge, and if it is much heavier than the standard Z.)

(3) The highest W_R^+ mass that can be attained is about 9 TeV, if values of σB (cross section times branching ratio) of 10^{-39} cm^2 can be studied in pp collisions. The heaviest Z_χ that could be seen under these conditions is about 6 TeV. For $\bar{p}p$ collisions with $\sigma B\geq 10^{-38}$ cm^2 the corresponding mass limits are $M_{W_R}^{\mp}\leq 7$ TeV, $M_{Z_\chi}\leq 4$ TeV. All these estimates are for $\sqrt{s}=40$ TeV.

(4) Distinctive aspects of W_R and Z_χ decays include (a) very hard leptons, whose nature and sign must be identified, (b) neutral leptons N which are probably heavy and which may travel a finite distance before decaying, and (c) characteristic forward-backward asymmetries of emitted leptons.

(5) Other heavy Z's also have been discussed. These include (a) a particle Z_ψ, coupled to the U(1) charge in $E_6>SO(10)\times U(1)_\psi$, (b) a Z_I^3, coupled to an SU(2) in $E_6\supset SU(6)\times SU(2)_I$, where the I (inert) SU(2) subgroup commutes with electric charge, and (c) a horizontal Z_H^8, member of an octet of SU(3), coupling differently to different generations. These illustrate the range of possibilities to which one should be receptive in multi-TeV collisions.

In what context are W_R's and Z's in the multi-TeV mass range interesting?

(1) Grand unified theories do not provide much guidance with regard to masses. There is some tendency for W_R to be extremely heavy ($\gtrsim 10^6$ GeV) in such theories.[43] The mass of the Z_χ is less constrained. However, the Z_χ cannot be too much lighter than N, which in turn is related to the W_R by a Yukawa coupling, or $N\bar{N}$ loops will contribute too much to the Z_χ mass. To summarize, the idea of W_R and Z_χ accessible at accelerator energies *may* be a long shot.

(2) Some hint of a possible multi-TeV scale for W_R comes from CP violation. If the scale for $m(K_S^0) - m(K_L^0) \equiv \Delta m$ sets the inequality $M_{W_R} \geq 1-2$ TeV,[10] and if the amplitude $\epsilon \simeq 2\times 10^{-3}$ for CP violation is at least in part due to W_R-W_L loops, we can estimate

$$M_{W_R} \sim \epsilon^{-1/2} \times (1-2\ \text{TeV}) = 20-50\ \text{TeV}\ . \tag{7.1}$$

This is above the values quoted in Table I, but not by much. A somewhat more optimistic estimate

$$M_{W_R} = 10\ \text{TeV} \times 2^{\pm 1} \tag{7.2}$$

is given by Harari and Leurer.[11] One can thus envision at least the *beginnings* of a direct assault on energy scales of interest in CP violation at the next generation of pp or $\bar{p}p$ colliders.

ACKNOWLEDGMENTS

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