Title
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Publication Date
1987-06-01

LBL-23694 c.2

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Presented at the Workshop on Physics at Future Accelerators, La Thuile, Val d'Aosta, Italy, January 7–13, 1987, and to be published in Volume II of the Proceedings

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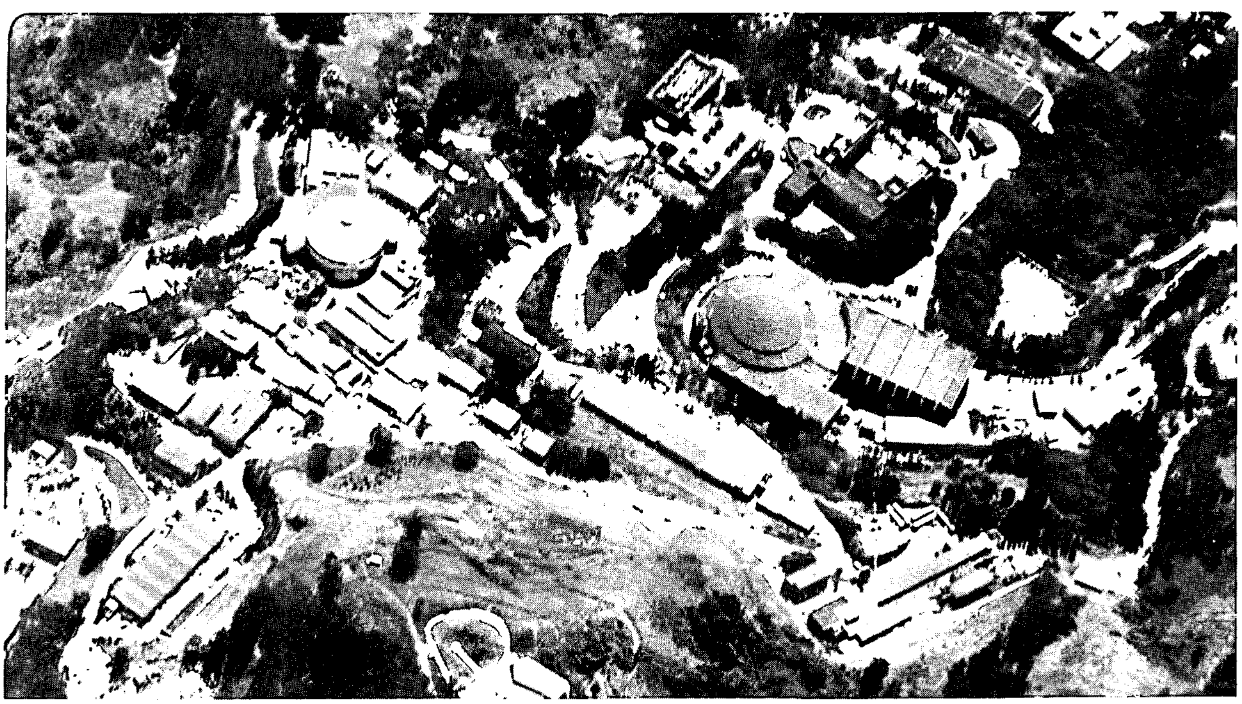
June 1987



Prepared for the U.S. Department of Energy under Contract DE-AC03-76SF00098

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LBL-23694

The physics of new Z' bosons[1]

F. del Aguila[2], M. Quirós[3] and F. Zwirner[4]

Abstract

We present a convenient parametrization for the couplings of a possible new neutral gauge boson (Z') with mass within the TeV region, and we exemplify it with three superstring-inspired models. We describe some general aspects of the Z' physics: mixing with the standard model Z^0 and its effects, present experimental limits, expected width and branching ratios, perspectives for Z' detection at future hadronic colliders.

(June 1987)

To appear in Vol. II of the Proceedings of the Workshop on Physics at Future Accelerators, La Thuile (Val d'Aosta) and CERN, 7-13 January 1987.

[1]This work was supported by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U. S. Department of Energy under Contract DE-AC03-76SF00098.

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The standard model of strong and electroweak interactions, based on the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$, is today in excellent agreement with all confirmed data in particle physics. However, there is no theoretical reason to believe that it must remain valid, in its present form, as soon as we gain experimental access to higher and higher energy scales. If there is some 'beyond the standard model' physics in the TeV region, which could be explored in the near future by some candidate accelerators (LHC, SSC, CLIC), an attractive possibility is another massive neutral gauge boson (Z') in addition to the Z discovered at the CERN $Sp\overline{p}S$ Collider[1)]. On the one hand, the limits which can be extracted from the present neutral currents data are not very stringent, and experimental searches at future accelerators seem to be relatively straightforward. On the other hand, the existence of additional neutral currents can be easily incorporated in non-minimal models of grand unification (e.g., those based on the groups $SO(10)$ or E_6), and this possibility naturally arises in the presently fashionable superstring models.

To set the framework for the study of a possible Z', one needs a convenient parametrization of its interactions: a general discussion, with references to previous work, can be found in Refs. 2-3. Here we assume that there is only one (elementary) extra Z' in the TeV region, and that it is associated to a flavour-conserving neutral current of E_6. Then the (fermionic) neutral current Lagrangian can be written as

$$\mathcal{L}_{NC} = \overline{\psi}^k \gamma_\mu (v_\alpha^k + a_\alpha^k \gamma_5) \psi^k Z_\alpha^\mu, \tag{1}$$

where summation over repeated indices is understood. The index $k = u, d, e, \ldots$ runs over the different physical fermions, while the index $\alpha = 1, 2, 3$ runs over the gauge boson mass eigenstates: $\alpha = 1$ corresponds to the photon $\gamma \equiv Z_1$, $\alpha = 2$ to the observed $Z \equiv Z_2$, $\alpha = 3$ to our hypothetical $Z' \equiv Z_3$. The explicit expressions for the vector and axial couplings of ordinary quarks and leptons ($u \equiv u, c, t$; $d \equiv d, s, b$; $e \equiv e, \mu, \tau$; $\nu \equiv \nu_e, \nu_\mu, \nu_\tau$) are shown in Table 1. A few explanatory comments are in order. To make the number of free parameters acceptably small, some more (plausible) assumptions have been made. (i) Fermions are supposed to lie in fundamental $\underline{27}$ representations of E_6, which contain, in addition to ordinary quarks and leptons, other exotic states: the main ambiguity in this embedding is parametrized by a mixing angle β, while other possible mixing angles are neglected. (ii) Charged currents are supposed to be described at low energy

$$v_2^u = \frac{e}{s_W c_W}\left[c_3\left(\frac{1}{4} - \frac{2}{3}s_W^2\right) - s_3\frac{5}{12}s_W s_1 c_1\left(\lambda - \frac{1}{\lambda}\right)\right]$$

$$v_3^u = \frac{e}{s_W c_W}\left[s_3\left(\frac{1}{4} - \frac{2}{3}s_W^2\right) + c_3\frac{5}{12}s_W s_1 c_1\left(\lambda - \frac{1}{\lambda}\right)\right]$$

$$a_2^u = \frac{e}{s_W c_W}\left\{-c_3\frac{1}{4} - s_3 s_W\frac{1}{\lambda}\left[\frac{1}{4}s_1 c_1\left(\lambda^2 - 1\right) + \frac{1}{3}\left(c_1^2 + \lambda^2 s_1^2\right)c_2\right]\right\}$$

$$a_3^u = \frac{e}{s_W c_W}\left\{-s_3\frac{1}{4} + c_3 s_W\frac{1}{\lambda}\left[\frac{1}{4}s_1 c_1\left(\lambda^2 - 1\right) + \frac{1}{3}\left(c_1^2 + \lambda^2 s_1^2\right)c_2\right]\right\}$$

$$v_2^d = \frac{e}{s_W c_W}\left\{c_3\left(-\frac{1}{4} + \frac{1}{3}s_W^2\right) + s_3\frac{1}{4}s_W\frac{1}{\lambda}\left[\frac{1}{3}s_1 c_1\left(\lambda^2 - 1\right) + \left(c_1^2 + \lambda^2 s_1^2\right)\left(c_2 + \sqrt{\frac{5}{3}}s_2\cos 2\beta\right)\right]\right\}$$

$$v_3^d = \frac{e}{s_W c_W}\left\{s_3\left(-\frac{1}{4} + \frac{1}{3}s_W^2\right) - c_3\frac{1}{4}s_W\frac{1}{\lambda}\left[\frac{1}{3}s_1 c_1\left(\lambda^2 - 1\right) + \left(c_1^2 + \lambda^2 s_1^2\right)\left(c_2 + \sqrt{\frac{5}{3}}s_2\cos 2\beta\right)\right]\right\}$$

$$a_2^d = \frac{e}{s_W c_W}\left\{c_3\frac{1}{4} + s_3\frac{1}{4}s_W\frac{1}{\lambda}\left[s_1 c_1\left(\lambda^2 - 1\right) + \left(c_1^2 + \lambda^2 s_1^2\right)\left(-\frac{1}{3}c_2 + \sqrt{\frac{5}{3}}s_2\cos 2\beta\right)\right]\right\}$$

$$a_3^d = \frac{e}{s_W c_W}\left\{s_3\frac{1}{4} - c_3\frac{1}{4}s_W\frac{1}{\lambda}\left[s_1 c_1\left(\lambda^2 - 1\right) + \left(c_1^2 + \lambda^2 s_1^2\right)\left(-\frac{1}{3}c_2 + \sqrt{\frac{5}{3}}s_2\cos 2\beta\right)\right]\right\}$$

$$v_\alpha^e = -2v_\alpha^u - v_\alpha^d \qquad (\alpha = 2,3)$$

$$a_\alpha^e = a_\alpha^d \qquad (\alpha = 2,3)$$

$$v_\alpha^\nu = -v_\alpha^d - \frac{1}{2}v_\alpha^u - \frac{1}{2}a_\alpha^u \qquad (\alpha = 2,3)$$

$$a_\alpha^\nu = -v_\alpha^\nu \qquad (\alpha = 2,3)$$

Table 1: Vector and axial couplings of $Z \equiv Z_2$ and $Z' \equiv Z_3$ to ordinary quarks and leptons.

by $SU(2)_L$, and only $SU(2)_L$ doublets and singlets are allowed to develop non-vanishing vacuum expectation values (VEVs). Under the previous assumptions, the unknown parameters which determine the Lagrangian of Eq. (1) are (apart from the usual ones of the standard model, β and the Z' mass) three angles (θ_1,θ_2,θ_3) and one ratio of coupling constants (λ). In our notation, $s_i \equiv \sin\theta_i$, $c_i \equiv \cos\theta_i$ ($i = 1,2,3,W$) and e is the running electromagnetic coupling constant. The angle θ_2 specifies the combination $\hat{Y}$ of E_6 generators which is broken only at low energy, in addition to the third component of the weak isospin, T_{3L}, and the usual weak hypercharge, Y. It is important to note that in superstring-inspired models, where the only gauge-non-singlet Higgs fields are contained in fundamental representations of E_6, θ_2 can take only two values: $\theta_2 = 0$ (corresponding to non-abelian Hosotani breaking at the compactification scale) and $\theta_2 = \arcsin\sqrt{\frac{5}{8}}$ (corresponding to large VEVs for some standard model singlets). Assuming grand unification, the parameter λ is generically of order 1, in the simplest models $\lambda = g_Y/g_{\hat{Y}}$. The angle θ_1 (usually ignored in the literature) describes the mixing of the currents associated to Y and $\hat{Y}$: except for special cases, a sizeable θ_1 is naturally generated by renormalization, since the $U(1)$ gauge couplings can mix already at the one loop level. The parameters λ and θ_1 are calculable in any specific model and their effects can be substantial[3–4], only in the particular case $\lambda = 1$ the angle θ_1 has no physical effect. Finally, θ_3 is the mixing angle characterizing the mass matrix for the (Z, Z') system.

To make the following discussion more concrete, we shall concentrate on three representative superstring-inspired models when giving numerical examples. The corresponding values of the parameters θ_2 and β are collected in Table 2. Model A corresponds to the case where E_6 is broken down to $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{\hat{Y}}$ directly at the compactification scale. Models B and C assume the existence of an intermediate stage of symmetry breaking, generated by large VEVs of some standard model singlets, and they differ only in the embedding of ordinary quarks and leptons inside the $\underline{27}$ of E_6. For simplicity, in all three models we have assumed $\lambda = 1$, so that physical quantities do not depend on θ_1.

The mass matrix for the (Z, Z') system can be written as

$$\mathcal{M}^2 = \begin{pmatrix} M^2_{ZZ} & M^2_{ZZ'} \\ M^2_{ZZ'} & M^2_{Z'Z'} \end{pmatrix} = \begin{pmatrix} c_3 & s_3 \\ -s_3 & c_3 \end{pmatrix} \begin{pmatrix} M^2_Z & 0 \\ 0 & M^2_{Z'} \end{pmatrix} \begin{pmatrix} c_3 & -s_3 \\ s_3 & c_3 \end{pmatrix}. \quad (2)$$

model	(c_2,s_2)	β
A	$(1,0)$	—
B	$(\sqrt{\frac{3}{8}},\sqrt{\frac{5}{8}})$	0
C	$(\sqrt{\frac{3}{8}},\sqrt{\frac{5}{8}})$	$\frac{\pi}{2}$

Table 2: Values of the parameters θ_2 and β corresponding to three representative 'superstring-inspired' models. We have assumed $\lambda = 1$, which implies no dependence on θ_1. There is no dependence on β in model A.

General expressions for M^2_{ZZ}, $M^2_{ZZ'}$ and $M^2_{Z'Z'}$ are given in Ref. 2-3. Here we limit ourselves to the three models introduced above, where only three VEVs contribute to the mass matrix of eq. (2), whose absolute values will be denoted by v, $\bar{v}$ and x. In our notation, v and $\bar{v}$ correspond to the $SU(2)_L$-doublet Higgses giving mass to u and (d,e), respectively, while x is associated to the standard model singlet which participates in the breaking of the residual $U(1)_{\tilde{Y}}$. In summary:

$$M^2_{ZZ} = \frac{e^2}{s_W^2 c_W^2}\frac{v^2+\bar{v}^2}{2} \equiv \frac{M_W^2}{c_W^2} = M^2_{Z^0} \quad (models\ A,B,C)\ ; \tag{3}$$

$$M^2_{ZZ'} = \begin{cases} \frac{e^2}{s_W c_W^2}\frac{\bar{v}^2-4v^2}{6} = \frac{s_W}{3}\frac{\bar{v}^2-4v^2}{v^2+\bar{v}^2}M^2_{Z^0} & (model\ A) \\ -\frac{e^2}{s_W c_W^2}\frac{v^2+\bar{v}^2}{\sqrt{6}} = -\sqrt{\frac{2}{3}}s_W M^2_{Z^0} & (model\ B) \\ \frac{e^2}{s_W c_W^2}\frac{3\bar{v}^2-2v^2}{2\sqrt{6}} = \frac{s_W}{\sqrt{6}}\frac{3\bar{v}^2-2v^2}{v^2+\bar{v}^2}M^2_{Z^0} & (model\ C) \end{cases}\ ; \tag{4}$$

$$M^2_{Z'Z'} = \begin{cases} \frac{e^2}{c_W^2}\frac{25x^2+16v^2+\bar{v}^2}{18} = \frac{s_W^2}{9}\frac{25x^2+16v^2+\bar{v}^2}{v^2+\bar{v}^2}M^2_{Z^0} & (model\ A) \\ \frac{e^2}{c_W^2}\frac{25x^2+4v^2+4\bar{v}^2}{12} = \frac{s_W^2}{6}\frac{25x^2+4v^2+4\bar{v}^2}{v^2+\bar{v}^2}M^2_{Z^0} & (model\ B) \\ \frac{e^2}{c_W^2}\frac{25x^2+4v^2+9\bar{v}^2}{12} = \frac{s_W^2}{6}\frac{25x^2+4v^2+9\bar{v}^2}{v^2+\bar{v}^2}M^2_{Z^0} & (model\ C) \end{cases}\ . \tag{5}$$

From Eqs. (2-5) one can extract some useful informations. To begin with, the entry M^2_{ZZ} is nothing else than the standard model prediction $M^2_{Z^0}$ for the Z^0 squared mass, which can be expressed as a function of the angle θ_W or of the mass M_W only. For $M^2_{Z'} > M^2_Z$, which is the case of interest here, a non-vanishing mixing term $M^2_{ZZ'}$ corresponds to an eigenvalue $M^2_Z < M^2_{Z^0}$, to be compared with the experimental value. The three parameters x, v and $\overline{v}$ in the mass matrix are therefore subject to two constraints: this severely restricts the possible values of the mixing angle θ_3. Note in particular that in model B $M^2_{ZZ'}$ is a function of θ_W only, implying $s_3 < 0$; also in model A negative values of s_3 are favoured if one accepts the theoretical indication that $\overline{v}/v \leq 1$. Only model C is naturally compatible with zero mixing.

Limits on the mass $M_{Z'}$ and on the mixing angle θ_3 can be obtained in each model by making a fit to the available neutral current data and to the measured values of the W and Z masses. At 90 % c.l. one obtains[5]:

$$Model\ A: \quad M_{Z'} > 130 GeV \quad -0.03 < \theta_3 < 0.20 \ \ (rad) \tag{6}$$

$$Model\ B: \quad M_{Z'} > 350 GeV \quad -0.01 < \theta_3 < 0.05 \ \ (rad) \tag{7}$$

$$Model\ C: \quad M_{Z'} > 180 GeV \quad -0.05 < \theta_3 < 0.03 \ \ (rad) \tag{8}$$

The relatively large positive values of θ_3 for model A are allowed only for Z' masses very close to the lower bound. For higher values of $m_{Z'}$, θ_3 rapidly decreases as $1/M^2_{Z'}$. The above limits have been derived assuming that only $SU(2)_L$ doublets and singlets contribute to the mass matrix $\mathcal{M}^2$, and they can be further improved by making use of the more model-dependent constraints coming from eqs. (4-5).

In order to estimate the Z' 'discovery limits' at future colliders, one needs information about its total width and branching ratios. Unfortunately, these quantities are strongly model-dependent, since the decay channels kinematically open to the Z' are function of the masses of the top quark, of the other exotic fermions in the $\underline{27}$ of E_6, and of possible supersymmetric particles. As a first approximation, one can neglect mass effects in the final state, $Z - Z'$ mixing and QCD corrections. In this case the ratio between the total width $\Gamma_{Z'}$ and the mass $M_{Z'}$ does not depend on the Z' mass and can be a useful reference number. For illustrative purposes, we consider here two extreme cases:

(I) Z' can decay only into the observed fermions and the top quark;

(II) Z' can decay into three families of 27 fermions and their supersymmetric partners.

model	case	$10^3 \times \Gamma_{Z'}/M_{Z'}$	$B(e^+e^-)\%$
A	(I)	6.5	3.6
	(II)	38.	0.6
B	(I)	12.	5.9
	(II)	38.	1.8
C	(I)	6.5	5.4
	(II)	38.	0.9

Table 3: Z' width and branching ratio into electron-positron pairs for three representative 'superstring-inspired' models.

Realistic values are likely to lie in between the predictions of cases (I) and (II). Case (I) must be corrected for the top quark mass, but also for the decays into essentially massless 'right-handed' neutrinos, which seem impossible to avoid in any of the three models. Case (II) must be corrected to take into account the masses of the exotic fermions and of the supersymmetric particles. Values of $\Gamma_{Z'}/M_{Z'}$ and of the branching ratio into electron-positron pairs are collected in Table 3. For the reader's convenience, we give below the general formulae for the partial widths of $Z \equiv Z_2$ and $Z' \equiv Z_3$, including mass and mixing effects but still neglecting QCD corrections:

$$\Gamma(Z_\alpha \to f\bar{f}) = \frac{M_\alpha}{12\pi}(1-4\eta_{f\alpha})^{1/2} C_f \left\{ (v_\alpha^f)^2 + (a_\alpha^f)^2 + 2[(v_\alpha^f)^2 - 2(a_\alpha^f)^2]\eta_{f\alpha} \right\}, \tag{9}$$

$$\Gamma(Z_\alpha \to \tilde{f}_{L,R}\bar{\tilde{f}}_{L,R}) = \frac{M_\alpha}{12\pi}(1-4\eta_{\tilde{f}\alpha})^{3/2} \frac{1}{4} C_f (v_\alpha^f \mp a_\alpha^f)^2, \tag{10}$$

In the above equations, v_α^f and a_α^f are the couplings of the generic (s)fermion f $(\tilde{f})$ to the gauge boson Z_α $(\alpha = 2,3)$, $\eta_{a\alpha} \equiv (M_a/M_\alpha)^2$ with $a = f, \tilde{f}$, and C_f is a colour factor equal to 1(3) for $SU(3)_C$ singlets (triplets). Obviously, only decays with $\eta_{a\alpha} < 1/4$ are kinematically allowed. If $M_{Z'} > 2M_W$ and there is non-zero $Z - Z'$ mixing, also the decay channel $Z' \to W^+W^-$ can be important. The corresponding decay rate is given by

$$\Gamma(Z' \to W^+W^-) = \frac{M_{Z'}}{12\pi}(1-4\eta_W)^{3/2}(1+20\eta_W+12\eta_W^2)\eta_W^{-2}\frac{\delta^2}{16}, \quad (11)$$

where $\delta = es_3c_W/s_W$ and $\eta_W \equiv (M_W/M_{Z'})^2$. An important point to note is that in the expression for the branching ratio $B(Z' \to W^+W^-) = \Gamma(Z' \to W^+W^-)/\Gamma_{Z'}$ the suppression factor $\sim (M_W/M_{Z'})^4$ coming from s_3^2 is compensated by the kinematical factor $\eta_W^{-2} = (M_{Z'}/M_W)^4$, so that one can have significant values for B even for very small values of θ_3. For example, in the three models under consideration $\Gamma(Z' \to W^+W^-)$ is generically of the same order of magnitude as $\Gamma(Z' \to e^+e^-)$. This does not hold, however, for supersymmetric decays into charged gauginos.

One way of detecting a Z' boson could be the observation of its indirect effects at future e^+e^- and ep colliders: these topics are discussed in other contributions to these Proceedings[6)]. The most attractive possibility, however, is the direct production of Z' bosons at present and future hadronic colliders: $Sp\bar{p}S$ (+$ACOL$), Tevatron, LHC and SSC.The most promising decay channel for the discovery of a Z' is, obviously, the one into charged lepton-antilepton pairs (electrons or muons), which led in the past to the discovery of the Z. Another interesting channel, which could give some useful complementary information, but is generally less favourable than the previous one, is the one into W^+W^- pairs. Detectability seems to rely on the chain $Z' \to W^+W^- \to l\ \nu\ jet\ jet$, but one has to find out a clever way of discriminating against the QCD background. Decays into ordinary quark-antiquark pairs appear to be undetectable, due to the huge QCD background. Other fancier possibilities require the presence of exotic fermions or supersymmetric particles in the final state, and will not be considered here.

The general form of the unpolarized differential cross-section for the process $p\bar{p} \to l^+l^-X$ ($l = e, \mu, \ldots$) is, for tree level amplitudes,

$$\frac{d\sigma}{dM\,dy\,d\cos\theta^*} = \Sigma_q[g_q^S(y,M)S_q(M)(1+\cos^2\theta^*) + g_q^A(y,M)A_q(M)2\cos\theta^*], \quad (12)$$

with

$$g_q^{S,A}(y,M) = \frac{M}{48\pi}x_a x_b[f_q^{(a)}(x_a,M^2)f_{\bar{q}}^{(b)}(x_b,M^2) \pm f_{\bar{q}}^{(a)}(x_a,M^2)f_q^{(b)}(x_b,M^2)], \quad (13)$$

where the sign $+(-)$ corresponds to $S(A)$, and

$$S_q = \Sigma_{\alpha,\beta} \frac{(v_\alpha^q v_\beta^q + a_\alpha^q a_\beta^q)(v_\alpha^l v_\beta^l + a_\alpha^l a_\beta^l)}{(M^2 - M_\alpha^2 + iM_\alpha\Gamma_\alpha)(M^2 - M_\beta^2 - iM_\beta\Gamma_\beta)}, \qquad (14)$$

$$A_q = \Sigma_{\alpha,\beta} \frac{(v_\alpha^q a_\beta^q + a_\alpha^q v_\beta^q)(v_\alpha^l a_\beta^l + a_\alpha^l v_\beta^l)}{(M^2 - M_\alpha^2 + iM_\alpha\Gamma_\alpha)(M^2 - M_\beta^2 - iM_\beta\Gamma_\beta)}. \qquad (15)$$

Quark and lepton masses are neglected. a and b are the two colliding hadrons at centre of mass energy $\sqrt{s}$, x_a and x_b being the momentum fractions of the colliding partons in a and b respectively. M is the lepton-antilepton invariant mass, y the rapidity ($x_{a,b} = M/e^{\pm y}\sqrt{s}$) and θ^* the scattering angle (al^-) in the centre of mass of the parton system $q\overline{q}$. Finally, $f_{q(\overline{q})}^{(a(b))}(x_{a(b)}, M^2)$ are the relevant parton distribution functions. Note that all the model-dependence is in S_q and A_q. Eqs. (12-15) are the starting point for the study of possible Z' signals at hadronic colliders. In Ref. 2 they have been used to compute cross sections and asymmetries for the different models at $Sp\overline{p}S$, Tevatron, LHC and SSC. The results of that analysis are the following. Present direct searches for a Z' at the CERN $Sp\overline{p}S$ still give weaker limits than the ones reported in eqs. (6-8), both for cases (I) and (II), but the situation can improve considerably when Tevatron will reach the design energy and luminosity. A more refined analysis is under way[7]. LHC and SSC should be able to probe Z' masses up to several TeV. More precise estimates of the discovery limits for the different models and machines have to take into account a number of additional factors. Experimentally, one needs to include cuts and detection efficiency[8], multiplying the cross section (12) by an experiment-dependent function of M, y and θ^*. Moreover, to measure asymmetries the lepton charges have to be identified and the momenta measured with a sufficient resolution: this could be non-trivial at LHC and SSC[9-10]. Theoretically, results depend on the parton distribution functions, different choices can correspond to slightly different results. More important, QCD corrections can also be non-negligible[10]. As a first indication, we give in the Figure the integrated cross section $\sigma(pp \to Z' \to e^+e^-) = \sigma(pp \to Z' \to \mu^+\mu^-)$ at LHC ($\sqrt{s} = 17\ TeV$) for the models and the cases considered above. Duke and Owens parton distribution functions[11] have been used. Further details and refinements can be found in Refs. 3,8,9,10.

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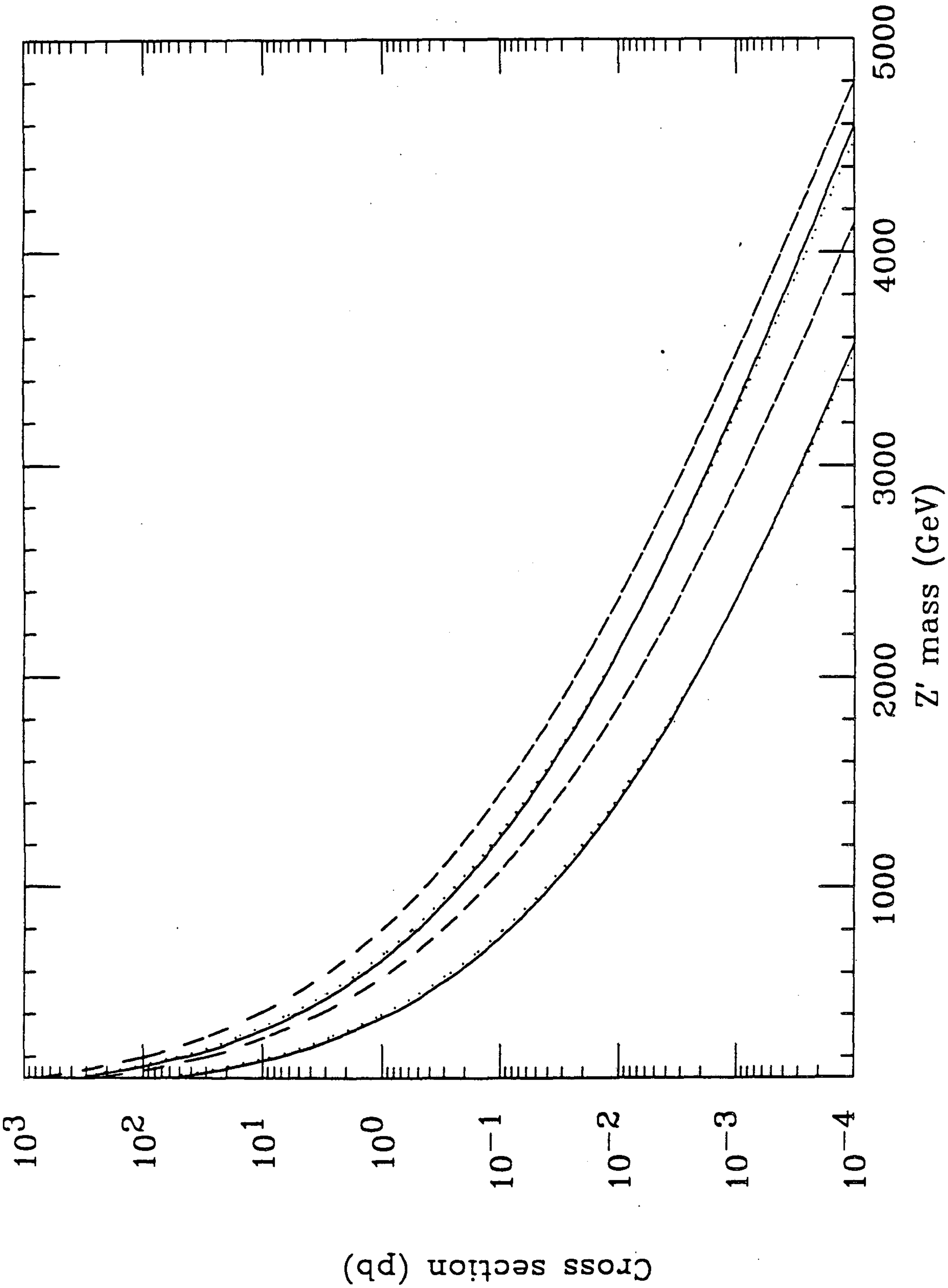
Figure caption:

Integrated cross section $\sigma(pp \to Z' \to e^+e^-) = \sigma(pp \to Z' \to \mu^+\mu^-)$ at the LHC ($\sqrt{s} = 17\ TeV$) as a function of the Z' mass. Solid lines correspond to model A, dashed lines to model B, dotted lines to model C. For each model, the upper and the lower lines correspond to cases (I) and (II), respectively.


Z' mass (GeV)
1000
2000
3000
4000
5000
Cross section (pb)
10^{3}
10^{2}
10^{1}
10^{0}
10^{-1}
10^{-2}
10^{-3}
10^{-4}

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