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NEW PHYSICS FROM PRECISION MEASUREMENTS

by

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New Physics from Precision Measurements

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Abstract

Precision tests of the electroweak theory in forthcoming experiments at the e^+e^- colliders LEP and SLC are discussed with respect to their sensitivity to new structures beyond the standard model. The signatures in the precisely measurable e^+e^- asymmetries on the Z resonance and the W mass are reviewed for various models of actual interest.

1 Introduction

The electroweak standard model is a very successful synthesis of a great deal of phenomena in the physics of elementary particles. At present there is no compelling experimental need for going beyond the minimal model. This success may encourage us to ask questions at a deeper level where the answers seem not to be within reach for the working period of the e^+e^- colliders LEP and SLC. This type of questions is tightly connected with our lack of understanding some basic ingredients of the model and gives naturally rise to the consideration of more complex structures than we encounter in the minimal model:

We do not understand the existence of families – why should the replication stop at three?

We do not understand the mechanism of electroweak symmetry breaking – why should the Higgs sector consist only of a single doublet?

We do not understand why the gauge group is SU(2)×U(1) – why should the underlying symmetry group not allow an extension SU(2)×U(1)×G at the Fermi scale with the existence of more (heavy) gauge bosons?

We do not understand the separation into bosons and fermions – why should the underlying structure not be supersymmetric but not obviously manifest at the presently accessible energy domain?

The list of possible extensions of the minimal model is by no means exhausted in terms of the scheduled questions. It might well be that there are alternative physical

Talk given at the Workshop on "Polarization with LEP", CERN, November 1987 systems that replace the job of the Higgs sector. If the symmetry breaking mechanism is due to a new strong force like in technicolor models, or if the vector bosons are composite objects, there may also be a new set of particles interacting by a new force described in terms of a mass scale in the TeV range. And, of course, there might be new dynamics which we have not yet thought about.

In a wider sense "new physics" may denote all types of empirically unknown physics in contrast to the "old physics" of the interaction between the vector bosons and the quarks and leptons. In that terminology also the bosonic sector of the minimal model including the Higgs particle would be a piece of new physics.

On the other hand it has meanwhile become a common language to classify as "new physics" everything which goes beyond the minimal model. We will in further course always refer to this more specific meaning when we are going to discuss possible effects of new physics in precision experiments at nowadays accessible energy domains.

From present days experiments there are no indications for structures beyond the minimal model. For this reason we may expect that any kind of new physics will manifest itself around the Z mass in terms of only small deviations from the standard model predictions rather than in overwhelming signatures, typically of the order of the conventional radiative corrections. The competition: standard radiative corrections – effects of new physics – requires a careful treatment of the radiative corrections in the minimal model which is at the moment complicated by the uncertainties induced by the unknown Higgs and the top quark mass.

The experimental verification of the standard loop effects will be a crucial step in confirming our confidence in the formulation of the electroweak theory as a quantized gauge theory. Since those effects which are not predominantly of QED origin are usually tiny (1% or less) the parameters of the minimal model have to be known with an adequate accuracy. To describe the interaction between fermions and gauge bosons we need essentially three parameters, e.g. the gauge coupling constants g_2, g_1, and the Higgs field vacuum expectation value v. Alternatively one may use [1-3]

$$\alpha,\ G_\mu,\ M_Z$$

(the electromagnetic fine structure constant, the Fermi constant, and the Z mass) which, together with the Z mass measurement in $e^+e^- \to f\bar{f}$ at LEP/SLC provide the input with the best accuracy (uncertainty $< 0.1\%$).

α and G_μ are experimental quantitites which are synonymous expressions for the Thomson cross section and the muon lifetime. Hence they are independent of any prejudice on the existence of an enriched particle scenario. M_Z will be measured from the shape of the resonance cross section, in particular from the location of the maximum. The relation between $\sqrt{s_{max}}$ and M_Z is influenced essentially by the initial state bremsstrahlung up to $O(\alpha^2)$ [4] and the s dependence of the Z width (ca. 35 MeV shift [5]). Both depend slightly (via the Z width) on the specific structure of the electroweak model. For an accuracy of about 20 MeV in the mass measurement, however, this small model dependence is not of practical importance, and M_Z may also be considered as an input parameter free of theoretical prejudice.

Besides the experiments defining the physical input one needs at least one additional experiment at the same level of accuracy for testing the underlying theory. Promising candidates are the W mass measurement and the polarization dependent asymmetries in $e^+e^- \to f\bar{f}$ which can be expressed in terms of the fermionic coupling constants v_f, a_f. Deviations from the tree level relations in the standard model

$$M_W^2\left(1-\frac{M_W^2}{M_Z^2}\right)=\frac{\pi\alpha}{\sqrt{2}G_\mu} \tag{1}$$
$$\frac{M_W^2}{M_Z^2\cos^2\theta_W}=1$$
$$v_f=\frac{I_3^f-2Q_f\sin^2\theta_W}{2\sin\theta_W\cos\theta_W},\quad a_f=\frac{I_3^f}{2\sin\theta_W\cos\theta_W}$$

can be interpreted in the following ways:

1. The deviations are due to the radiative corrections in the minimal model. Then it must be possible to settle a domain of Higgs and top masses M_H, m_t (favored within $10 < M_H < 1000$ GeV, $50 < m_t < 200$ GeV) in order to reproduce the experimental results at the level of radiative corrections .

2. The deviations are due to new particles which contribute to the relations between the measurable quantities through their virtual effects in the loops, but not at the tree level. This class would enclose new fermion generations, more Higgs doublets, the minimal supersymmetric extension of the standard model , and preserve in particular the relation between the vector boson masses and the mixing angle in (1). In order to be considered a convincing signal the observed effects should lie outside the predictions of the parameter range covered by 1.

3. The deviations are due to new structures which need to modify the standard model already at the tree level. The modification
$$\cos^2\theta_W=\frac{M_W^2}{\rho M_Z^2}$$
with an additional parameter ρ would become necessary e.g. in the presence of Higgs representations with dimension ≥ 3. New Z bosons give rise to a difference in the $\sin^2\theta_W$ dependence of the coupling constants v_f, a_f. A signal for the presence of this class of new contributions would be that the value for $\sin^2\theta_W$ as derived from the experimental vector boson masses: $\sin^2\theta_W = 1 - M_W^2/M_Z^2$, disagrees with the value obtained from asymmetry measurements, i.e. from the v_f/a_f ratio.

According to the various possibilities outlined above we subdivide our discussion as follows: First we want to give a brief repetition of the structure of the radiative corrections in the minimal model, in particular with respect to those experimental quantities which fulfill the requirements on accuracy. A general overview over current extensions of the minimal model follows and an inventory of existing work is given. Examples and numerical results of new physics effects are then listed for concrete situations of new generations of fermions, more Higgs doublets, the supersymmetric extension of the standard model , new Z bosons, and new vector bosons in composite models or models with a strong interacting Higgs sector. Finally we are concerned with the question of separating effects from new bosons not within SU(2)×U(1) from those induced by new (virtual) particles within the minimal gauge group. The present contribution may be understood as an update of [6] where already a lot of results for various kinds of new physics effects has been collected.

2 Precision experiments and the Standard Model

2.1 Synopsis of radiative corrections

For an adequate theoretical discussion aiming the accuracy of the high precision experiments at LEP/SLC the precise knowledge of the standard model radiative corrections becomes indispensable. During the last years it has become customary to work in the electroweak on-shell scheme (see e.g. [1-3] and references therein) for the calculation of higher order effects which is based on the parameters

$$\alpha, \ M_W, \ M_Z, \ M_H, \ m_f \tag{2}$$

together with the definition of the mixing angle as proposed by Sirlin [7]

$$\sin^2\theta_W = 1 - \frac{M_W^2}{M_Z^2}\,. \tag{3}$$

It has also become customary, and in the on-shell scheme quite naturally, to subdivide the one-loop corrections to e^+e^- processes in the following way [3] :

- QED corrections, which consist of those diagrams with an extra photon added to the Born diagrams either as a real bremsstrahlung photon or a virtual photon loop (Fig. 1). This class forms a gauge invariant subset and allows a separate treatment. Although considered not very interesting with respect to the underlying theory they are in general large at LEP energies and hence need a lot of attention for practical purposes.

- Weak corrections, which collect all other one-loop diagrams. The subset of diagrams which involve corrections to the vector boson propagators (γ, Z, W) (Fig. 2) are frequently called "oblique corrections" [6], in order to delimit them from the residual set of vertex corrections and box diagrams of Fig. 3 ("direct corrections" in the terminology of [6]).

Since Higgs boson contributions in Fig. 3 can be neglected for $m_f^2 \ll M_W^2$ because of small Yukawa couplings only Z and W bosons are present in virtual states. This is different for the propagator corrections which involve all particles of the model, in particular the as yet unknown Higgs particle and the top quark. Hence they depend

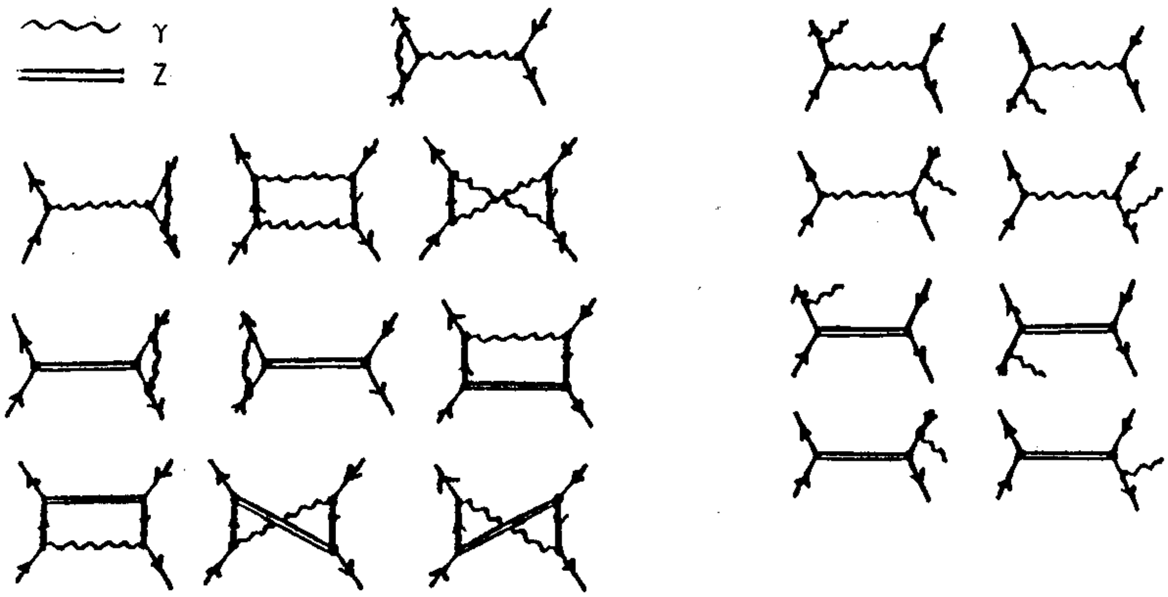


Figure 1: QED corrections for $e^+e^- \to f\bar{f}$

Figure 2: Propagator corrections

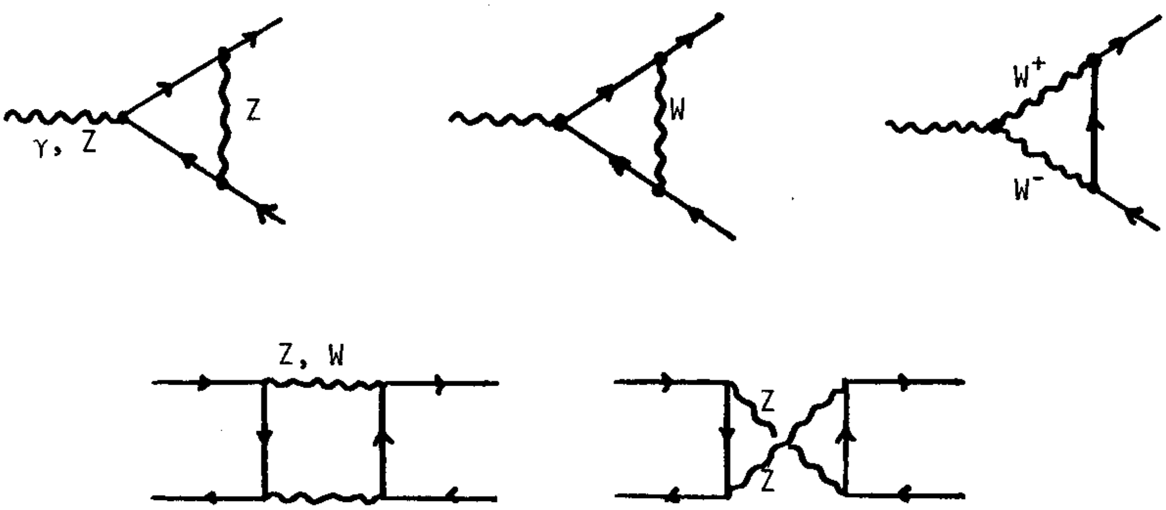


Figure 3: Vertex corrections and box diagrams (direct corrections)

on M_H and m_t, presently a source of uncertainty. Once calculated in terms of 1-particle irreducible renormalized self energies $\hat{\Sigma}$ the propagator corrections can easily be implemented in an effective matrix element of the Born type by replacing the lowest order propagators by dressed propagators D^{γ}, D^{Z} inverting the matrix

$$\begin{pmatrix} q^2 + \hat{\Sigma}^{\gamma}(q^2) & \hat{\Sigma}^{\gamma Z}(q^2) \\ \hat{\Sigma}^{\gamma Z}(q^2) & q^2 - M_Z^2 + \hat{\Sigma}^{Z}(q^2) \end{pmatrix} \tag{4}$$

and

$$D^W = \frac{1}{q^2 - M_W^2 + \hat{\Sigma}^W(q^2)} \,. \tag{5}$$

Due to the inversion of the matrix (4) also higher order terms are present in the propagators which become of some importance if e.g. the top quark is very heavy. For a simplified discussion we can approximate the neutral boson propagators by

$$D^{\gamma} = \left(q^2 + \hat{\Sigma}^{\gamma}(q^2)\right)^{-1} \tag{6}$$

$$D^{Z} = \left(q^2 - M_Z^2 + \hat{\Sigma}^{Z}(q^2)\right)^{-1} \,. \tag{7}$$

The $\gamma - Z$ mixing can be interpreted as a redefinition of the vector coupling constants (1) by

$$v_f \;\rightarrow\; v_f - Q_f \cdot \Pi^{\gamma Z} \tag{8}$$

with

$$\Pi^{\gamma Z} = \frac{\hat{\Sigma}^{\gamma Z}(q^2)}{q^2} \,. \tag{9}$$

Any kind of extra particles with couplings to the gauge bosons which do not couple directly to the external fermions would enter via the loop contributions to these 2-point functions.

The vertex corrections can be represented in terms of form factors $F_{V,A}^{Z(\gamma)f}(q^2)$. Also the box diagrams can be written as

(initial current) · (final current) · (formfactor).

The currents have only vector and axial vector contributions. All the formfactors can be found in analytical form in [8].

The relation $M_W \leftrightarrow M_Z \leftrightarrow \sin^2\theta_W$:

Besides the list of loop contributions to e^+e^- processes the radiative corrections to the charged current process

$$\mu^- \rightarrow e^- \, \nu_\mu \, \bar{\nu}_e$$

play a central role in the discussion of the electroweak parameters. The reason is the precisely measured Fermi constant G_μ which is one of the unambigously known input parameters. The μ lifetime τ_μ resp. G_μ establishes a relation between M_W and M_Z (as well as $\sin^2\theta_W$ by means of (3)). At the tree level this relation is given by the first equation of (1).

The one-loop (non-QED) corrections to τ_μ yield the following modification of (1) [7,9]:

$$M_W^2\left(1-\frac{M_W^2}{M_Z^2}\right)=\frac{\pi\alpha}{\sqrt{2}G_\mu}\cdot\frac{1}{1-\Delta r(\alpha,M_W,M_Z,M_H,m_t)}\,. \tag{10}$$

The radiative correction Δr is the sum of the W self energy at $q^2\approx 0$,which contains the Higgs and top dependence, and vertex and box diagrams involving only W and Z:

$$\Delta r=\frac{\hat{\Sigma}^W(0)}{M_W^2}+\frac{\alpha}{4\pi\sin^2\theta_W}\left(6+\frac{7-4\sin^2\theta_W}{2\sin^2\theta_W}\log\cos^2\theta_W\right)\,. \tag{11}$$

The renormalized W self energy can be expressed in terms of the unrenormalized self energies Σ^γ, Σ^Z, Σ^W, $\Sigma^{\gamma Z}$ which are the sum of all contributing one-loop diagrams (fermions, gauge bosons and ghosts, Higgs and Goldstone Higgs) in the following way:

$$\begin{aligned}\frac{\hat{\Sigma}^W(0)}{M_W^2)} &= \frac{\Sigma^W(0)-\Sigma^W(M_W^2)}{M_W^2}\\ &+\Pi^\gamma(0)+2\frac{\cos\theta_W}{\sin\theta_W}\frac{\Sigma^{\gamma Z}(0)}{M_Z^2}\\ &-\frac{\cos^2\theta_W}{\sin^2\theta_W}Re\left(\frac{\Sigma^Z(M_Z^2)}{M_Z^2)}-\frac{\Sigma^W(M_W^2)}{M_W^2}\right)\end{aligned} \tag{12}$$

where

$$\Pi^\gamma(0)=\frac{\partial\Sigma^\gamma}{\partial q^2}(q^2=0)\,. \tag{13}$$

The relation (10) is of twofold importance:

– It allows a a comparison of the theoretical $M_W - M_Z - \sin^2\theta_W$ correlation with the values obtained from precision experiments.

– It provides a value for M_W (after specifying the other masses of the model) as well as for $\sin^2\theta_W$, which can be used as input parameters for the calculation of v_f, a_f in (1) and hence for the observables in $e^+e^-\to f\bar{f}$ experiments like cross sections and asymmetries .

Fixing the Z mass at $M_Z = 92$ GeV the dependence of Δr on the Higgs and top mass, as derived from (10), can be read off from Figure 4.

The uncertainty in Δr associated with the error in the hadronic contribution to the photon vacuum polarization [10]

$$\delta(\Delta r)=\pm 0.0007$$

induces an uncertainty in M_W

$$\frac{\delta M_W}{M_W}=\frac{\sin^2\theta_W}{\cos^2\theta_W-\sin^2\theta_W}\cdot\frac{\delta(\Delta r)}{2(1-\Delta r)}$$

of about 12 MeV; together with the error following from the scheme dependence [11] the total uncertainty

$$\delta(\Delta r)=\pm 0.0013$$

yields $\delta M_W \approx 20$ MeV, which is not of practical importance in view of the accuracy $(\Delta M_W)_{exp} = 100$ MeV in the W mass measurement of the future [12].

A translation of $\delta(\Delta r)$ into the uncertainty for $\sin^2\theta_W$ yields with $\delta(\Delta r) = \pm 0.0013$:

$$\delta \sin^2\theta_W = \frac{\sin^2\theta_W \cos^2\theta_W}{\cos^2\theta_W - \sin^2\theta_W} \cdot \frac{\delta(\Delta r)}{1-\Delta r} \approx \pm 0.0004 \, .$$

This theoretical error matches the experimental accuracy with which $\sin^2\theta_W$ can be measured from polarized e^+e^- annihilation on the Z resonance.

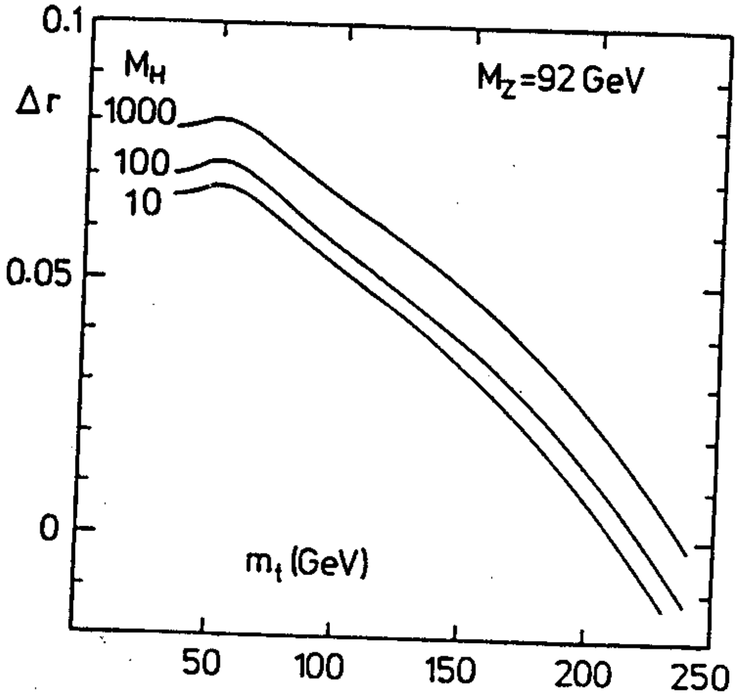


Figure 4: m_t and M_H dependence of Δr

2.2 On-resonance observables

We will assume that the Z mass can be measured from the resonance shape within 50 MeV (eventually 20 MeV). According to our general strategy this completes our set of input parameters α, G_μ, M_Z for fixing the theory.

For testing the theory further precisely measurable quantities are necessary. The requirement for any observable probing the electroweak structure is to be as insensitive as possible to theoretical and experimental systematic uncertainties. Besides the error associated with the hadronic input in the vector boson 2-point functions we have to consider additional theoretical uncertainties: the unknown effects of (possibly large) higher order ($\geq O(\alpha^2)$) QED contributions and the uncertainty in the QCD corrections for hadronic final states from the error in α_s. The statistical error is minimized near the top of Z resonance.

From the Z line shape one can determine, besides M_Z, the following quantities:

- the total cross section

$$\sigma_T = \sum_f \sigma(e^+e^- \to f\bar{f})$$

- the total Z width Γ_Z .

Without beam polarization the further observables are accessible:

- the partial widths $\Gamma_Z(Z \to f\bar{f})$;
- the forward-backward asymmetries

$$A_{FB}(e^+e^- \to f\bar{f}) = \frac{\sigma^{for} - \sigma^{back}}{\sigma^{for} + \sigma^{back}} \; ; \tag{14}$$

with

$$\sigma^{for} = \int_{\theta<\pi/2} d\Omega \, \frac{d\sigma}{d\Omega}, \quad \sigma^{back} = \int_{\theta>\pi/2} d\Omega \, \frac{d\sigma}{d\Omega}$$

The instrument of longitudinal polarization for the incident e^- beam offers the polarization dependent asymmetries

- the left-right asymmetry

$$A_{LR} = \frac{\sigma_L - \sigma_R}{\sigma_L + \sigma_R} \tag{15}$$

where $\sigma_{L(R)}$ denotes the integrated cross section for left (right) handed electrons;

- the polarized forward-backward asymmetry

$$A^{pol}_{FB}(e^+e^- \to f\bar{f}) = \frac{\sigma_L^{for} - \sigma_L^{back} - (\sigma_R^{for} - \sigma_R^{back})}{\sigma_L^{for} + \sigma_L^{back} + \sigma_R^{for} + \sigma_R^{back}} \tag{16}$$

In lowest order the on-resonance asymmetries are simple combinations of the fermionic coupling constants [1] :

$$A_{FB} = \frac{3}{4} \cdot \frac{2v_e a_e}{v_e^2 + a_e^2} \cdot \frac{2v_f a_f}{v_f^2 + a_f^2} \tag{17}$$

$$A_{LR} = \frac{2v_e a_e}{v_e^2 + a_e^2} \tag{18}$$

$$A_{FB}^{pol} = \frac{3}{4} \cdot \frac{2v_f a_f}{v_f^2 + a_f^2} \tag{19}$$

The ingredients in all these asymmetries are the combinations

$$A_f := \frac{2v_f a_f}{v_f^2 + a_f^2} = \frac{2 \mid 2I_3^f - 4Q_f \sin^2\theta_W \mid}{1 + (2I_3^f - 4Q_f \sin^2\theta_W)^2} \tag{20}$$

which are functions of $\sin^2\theta_W$ only for a given species of fermions. Note that $A_{FB}^{pol}(e^+e^- \to f\bar{f})$ measures the same quantity A_f as the final state polarization of the fermion f in $e^+e^- \to f\bar{f}$.

Since a sensitivity to $\sin^2\theta_W$ means also a sensitivity to the internal structure of the weak part of the model (because of the correlation (10)) the structure of (17-19) suggests A_{LR} and A_{FB}^{pol} to be the best candidates for electroweak tests: A_{LR} measures the initial state and A_{FB}^{pol} the final state coupling constants, allowing simultanously tests of the fermion universality. The particular sensitivity of A_{LR} to $\sin^2\theta_W$ will allow a measurement of this parameter with an error of ± 0.0004 if an accuracy $(\Delta A_{LR})_{exp} = 0.003$ can be achieved (Fig. 5).

An important advantage of A_{LR} is the property that it is only very little influenced by the (otherwise large) QED corrections [13] and by QCD corrections in case of quark pair production [14] as far as the quark masses are negligible. Moreover, also A_{FB}^{pol} has been shown to get only very small QED corrections [15], allowing the conjecture that the next order contributions do not play a practical role.

A quantity of comparable interest and with similar features is the τ polarization in $e^+e^- \to \tau^+\tau^-$, which, in case of lepton universality, fulfills

$$P_\tau = A_{LR} = \frac{4}{3} \cdot A_{FB}^{pol}(e^+e^- \to \mu^-\mu^+)\ . \tag{21}$$

This remains valid also when radiative corrections are included in the minimal model (for a more detailed discussion see the contribution by Wąs [16]).

The sensitivity of A_{LR} and M_W with respect to m_t and M_H is illustrated in Figure 6. Eliminating $\sin^2\theta_W$ by means of (10) leads to the prediction (for fixed M_Z)

$$M_Z,\ G_\mu,\ \alpha \longrightarrow M_W(M_H, m_t),\ A_{LR}(M_H, m_t)$$

The curves for fixed m_t (dashed lines) and for fixed M_H (full lines) are displayed in Figure 6. The weak corrections coming from the top and the Higgs, and the expected

[1] up to small terms $\sim (\Gamma_Z/M_Z)^2$ coming from the pure γ exchange

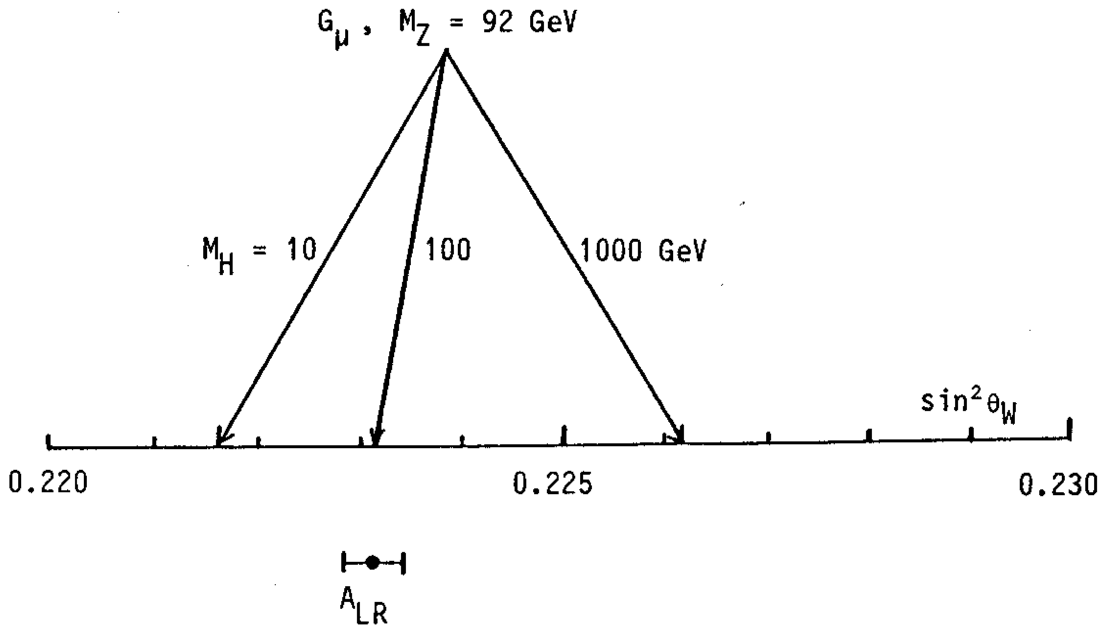


Figure 5: $\sin^2\theta_W$ from α, G_μ, M_Z

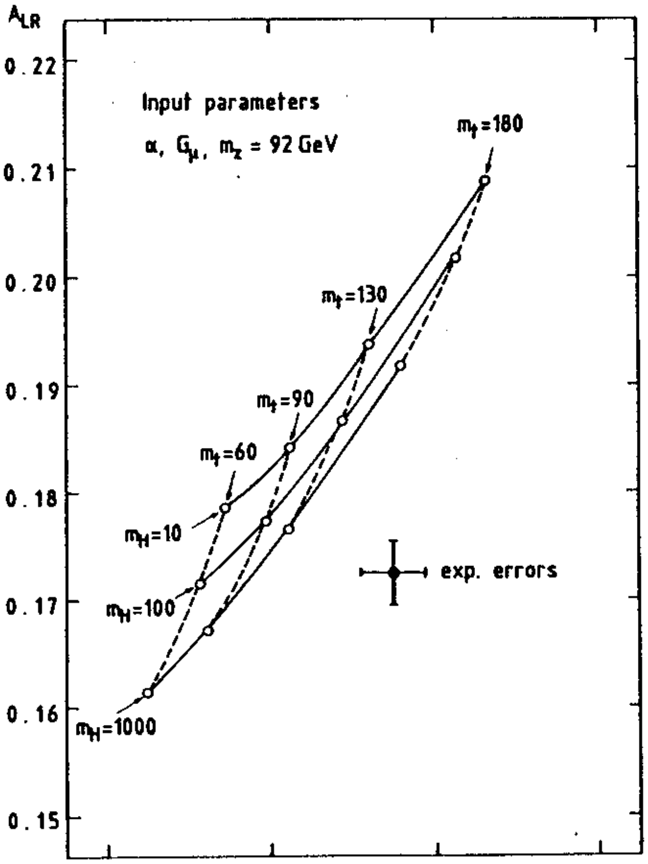


Fig. 6. Variation of the predictions for the left-right asymmetry A_{LR}, and for m_W, as a function of m_t and m_H, in comparison with the expected experimental precision of A_{LR} and m_W. (From /3/, based on the results in /2/.)

experimental errors make it possible to delimit the allowed mass ranges significantly. Effects from new physics can be revealed if they give rise to values of M_W and A_{LR} outside the domain covered in Figure 6.

3 Inventory of new physics effects

The structure of the standard model makes it possible to calculate the predictions for all measurable quantities with the required accuracy. Considering the situation possibly beyond the minimal model we do not know which type of new physics would become visible. But without being apparently conservative we may assume that it will show up as small deviations suggesting that a perturbative treatnment is allowed. In order to decide whether a particular question can be answered at the given facilities we have to work out the consequences of any particular hypothesis. To this aim we want to specify more concretely the list of possible extensions of the standard model which are of current interest, and to classify them according to our strategy outlined in the introduction. Table 1 schedules a list of new physics of current interest and the way how the measurable quantities resp. the theoretical relations between our observables are influenced.

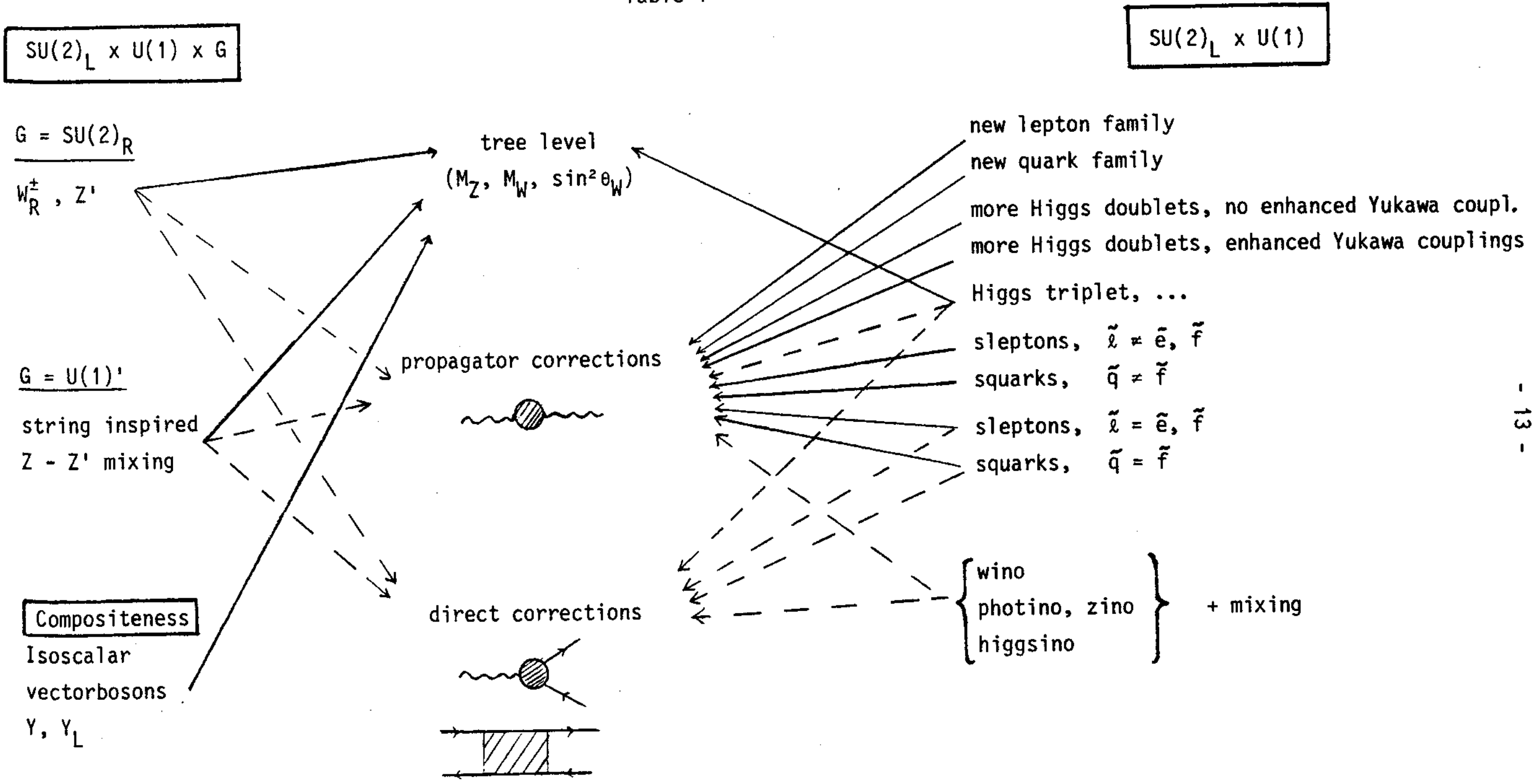
The relevant relations are those between the boson masses M_W, M_Z, and $\sin^2\theta_W$ as measured from the on-resonance asymmetries. New (heavy) particles will either contribute via their presence in the loops only or they will already become manifest in the tree level relations.

The first situation corresponds to any type of new structure that can be embedded in the minimal gauge group leaving the fundamental relations (1) in lowest order invariant. The further subdivision into objects entering only the gauge boson propagators ("oblique") and those with "direct" couplings to the fermions in $e^+e^- \to f\bar{f}$ resp. μ decay is for practical reasons mainly: The "oblique" contributions can be treated in a systematic way without specifying the detailed content of the model. Several examples have already been discussed in the literature [6]. For objects with direct couplings to the external fermions more detailed information is needed and hence more model dependence is introduced.

The second possibility: existence of new structures which modify (some of) the relations (1) already at the tree level, requires an extension of the minimal parameter set (2) by at least one additional quantity, e.g. the mixing angle as an independent parameter. The presence of new currents as in left-right symmetric models [17] or in models inspired by superstrings [18] distorts the dependence of the coupling constants in (1) on only $\sin^2\theta_W$ by the existence of an additional mixing angle. The character of this class of new physics is more complicated since the new objects not only contribute in lowest order but augment the radiative corrections sector as well. Whereas an extensive literature on lowest order effects has already been accumulated [17-20] very little has been performed yet at the level of complete radiative corrections [20].

The signature of the arrows in table 1 is as follows: the arrows indicate how the corresponding physical object finds its way into the relations between the measurable

Table 1



$SU(2)_L \times U(1) \times G$
$G = SU(2)_R$
$W_R^{\pm}$, Z'
$G = U(1)'$
string inspired
Z - Z' mixing
Compositeness
Isoscalar
vectorbosons
Y, Y_L
tree level
(M_Z, M_W, $\sin^2\theta_W$)
propagator corrections
direct corrections
$SU(2)_L \times U(1)$
new lepton family
new quark family
more Higgs doublets, no enhanced Yukawa coupl.
more Higgs doublets, enhanced Yukawa couplings
Higgs triplet, ...
sleptons, $\tilde{\ell} \neq \tilde{e}, \tilde{f}$
squarks, $\tilde{q} \neq \tilde{f}$
sleptons, $\tilde{\ell} = \tilde{e}, \tilde{f}$
squarks, $\tilde{q} = \tilde{f}$
wino
photino, zino
higgsino
+ mixing

quantities. Full arrows denote thereby that the effect has been calculated and discussed in the currently actual models; dashed arrows signalize that the corresponding path has either not yet or only insufficiently been treated at the present stage. The radiative corrections are then customarily incorporated by taking over the standard model set (eventually with slight modifications).

4 New particles in SU(2)×U(1) preserving the tree level structure in 4-fermion processes

4.1 Systematic aspects

Many physically intriguing extensions of the standard model exhibit a structure obeying the minimal gauge group SU(2)×U(1) and preserving the tree level formulae for fermionic processes of the standard model . Typical examples of this class are extra fermion families, more Higgs doublets, scalar fermions, ...

Preserving the tree level relations means more precisely:

That part of the Lagrangian of the extended model which describes the 4-fermion processes in lowest order is identical to the corresponding part of the standard model Lagrangian. This property induces a structure in the renormalized one-loop radiative corrections which is formally completely equivalent to that of the minimal radiative corrections. The structure of the counter terms, the number of the renormalization constants, and the renormalization conditions can be taken over from the standard model as formulated e.g. in the on-shell scheme [8]. The only difference is the appearance of extra loop diagrams with the new particles giving rise to additional terms in the gauge boson self energies and, in case, in the vector and axial vector form factors for vertex and box contributions.

The situation becomes particularly simple if the new particles have no direct couplings to the fermions. The following set of formulae is as general as to comprise all specific situations of these "oblique" corrections [2]. Since the only place where such new objects can enter the radiative corrections are the gauge boson propagators (5-9) it is sufficient to deal with the renormalized vector boson self energies. Expressed in terms of the corresponding sums of all contributing diagrams they are given by

$$
\begin{aligned}
\hat{\Sigma}^{\gamma}(q^2) &= \Sigma^{\gamma}(q^2) - \Pi^{\gamma}(0)\cdot q^2 , \qquad (22)\\
\hat{\Sigma}^{\gamma Z}(q^2) &= \Sigma^{\gamma Z}(q^2) - \Sigma^{\gamma Z}(0) + q^2 \left\{ 2\frac{\Sigma^{\gamma Z}(0)}{M_Z^2} - \frac{\cos\theta_W}{\sin\theta_W}\left(\frac{\delta M_Z^2}{M_Z^2} - \frac{\delta M_W^2}{M_W^2}\right)\right\} ,\\
\hat{\Sigma}^{Z}(q^2) &= \Sigma^{Z}(q^2) - \delta M_Z^2 + \delta Z_2^Z (q^2 - M_Z^2),\\
\hat{\Sigma}^{W}(q^2) &= \Sigma^{W}(q^2) - \delta M_W^2 + \delta Z_2^W (q^2 - M_W^2)
\end{aligned}
$$

[2] A discussion equivalent to ours is also given in [6]

with $\Pi^\gamma(0)$ from (13) and

$$\begin{aligned}
\delta Z_2^Z &= -\Pi^\gamma(0) - 2\,\frac{\cos^2\theta_W - \sin^2\theta_W}{\sin\theta_W\cos\theta_W}\,\frac{\Sigma^{\gamma Z}(0)}{M_Z^2} \\
&\quad + \frac{\cos^2\theta_W - \sin^2\theta_W}{\sin^2\theta_W}\left(\frac{\delta M_Z^2}{M_Z^2} - \frac{\delta M_W^2}{M_W^2}\right) \qquad (23)\\
\delta Z_2^W &= -\Pi^\gamma(0) - 2\,\frac{\cos\theta_W}{\sin\theta_W}\,\frac{\Sigma^{\gamma Z}(0)}{M_Z^2} + \frac{\cos^2\theta_W}{\sin^2\theta_W}\left(\frac{\delta M_Z^2}{M_Z^2} - \frac{\delta M_W^2}{M_W^2}\right) \\
\delta M_{Z,W}^2 &= Re\,\Sigma^{Z,W}(M_{Z,W}^2)\,.
\end{aligned}$$

After specification of the model dependent coupling constants and the masses all quantities on the r.h.s of (22) fixed; the results are finite and expressed in terms of our basic parameter set (2) augmented by the parameters of the extra particles.

The mixing angle $\sin^2\theta_W$ as well as M_W are related to M_Z formally in the same way as in the standard model by eq. (10). The presence of new particles is reflected in an additive term in Δr

$$\Delta r = (\Delta r)_{min} + \left(\frac{\hat{\Sigma}^W(0)}{M_W^2}\right)_{new} \qquad (24)$$

which can be calculated by use of (12).

The on-resonance asymmetries (17-19) also allow a simple discussion: Since they are (up to small terms) determined by the quantity A_f, eq.(20) and therefore by the ratio v_f/a_f it is sufficient to study the response of this ratio to the presence of the new particles.

Because of the modification (8) of the vector coupling by the $\gamma - Z$ mixing we have

$$\begin{aligned}
\frac{v_f}{a_f} &\rightarrow \frac{I_3^f - 2\,\sin^2\theta_W\,Q_f\left(1 + \frac{\cos\theta_W}{\sin\theta_W}\,\Pi^{\gamma Z}(M_Z^2)\right)}{I_3^f} \\
&= \frac{I_3^f - 2\,\overline{\sin}^2\theta_W\,Q_f}{I_3^f} \qquad (25)
\end{aligned}$$

with an effective $\overline{\sin}^2\theta_W$ which has the same value for all types of fermions. Hence all the on-resonance asymmetries can still be parametrized with help of a single quantity. A verification of this universality of $\overline{\sin}^2\theta_W$ in the various asymmetries is therefore a helpful instrument in the systematic analysis of deviations from the standard model. This will become of particular importance in isolating effects of extra Z bosons, as discussed in Section 6.

The presence of additional direct couplings to fermions will spoil this simple pattern since the induced form factors become fermion type dependent and the results can in general not be written in a model independent way. For many discussions these corrections have been neglected because they are usually small as we know from the standard model. In high precision experiments, however, they can reach the level of experimental observability.

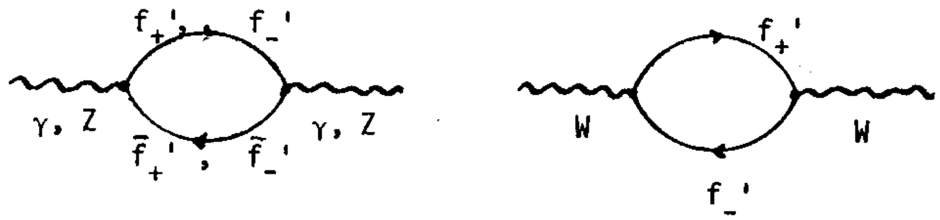
In the following we want to discuss as concrete examples:
– effects from an extra quark/lepton family;
– effects from scalar fermions;
– effects from a second Higgs doublet.
Part of the results put together here are already contained in [6].

4.2 New lepton and quark families

The presence of a new fermion generation $\begin{pmatrix} f'_+ \\ f'_- \end{pmatrix}$ leads to a repetition of the pattern in the known fermion contributions. The essential feature is the existence of large radiative corrections in case of a big mass difference between the isospin states of the doublet [21]. The contribution to Δr

$$\Delta r = \frac{\alpha}{4\pi} \cdot \frac{3\sin^2\theta_W}{4\cos^2\theta_W} \cdot \frac{m'^2_+ - m'^2_-}{M_W^2} + \cdots$$

from the diagrams



cause a shift in the W mass and also, via the induced shift in $\sin^2\theta_W$ as well as the $\gamma - Z$ mixing term in (25), a shift in the sensitive asymmetries determined by A_f, eq. (20).

The shift δA_{LR} in the left-right asymmetry due to a new quark doublet is shown in Figure 7a for various mass ratios of the up-down members and analogously in Figure 7b for a new lepton doublet, together with the experimental accuracy. The sensitivity to mass splittings is obvious. It should be noticed (as emphasized already in [6]) that the presence of a new family will give rise to a small shift in A_{LR} also in the case of mass degeneracy. This shift is independent of the heavy mass scale; for one extra generation it is of the order of the experimental error.

An important question is the separability of the predicted deviations from the domain covered by the standard model . This will be elucidated in Figure 8 where, in a plot similar to Figure 6, the simultanous effects in M_W and A_{LR} are displayed for new leptons and quarks together with the standard model model curves induced by the Higgs uncertainty. From an optimistic point of view a separability seems possible for a wide range of parameters. However, since a standard heavy top quark has a similar signature as the extra doublet the effects will be obscured as far as m_t is the great unknown quantity in the minimal model.

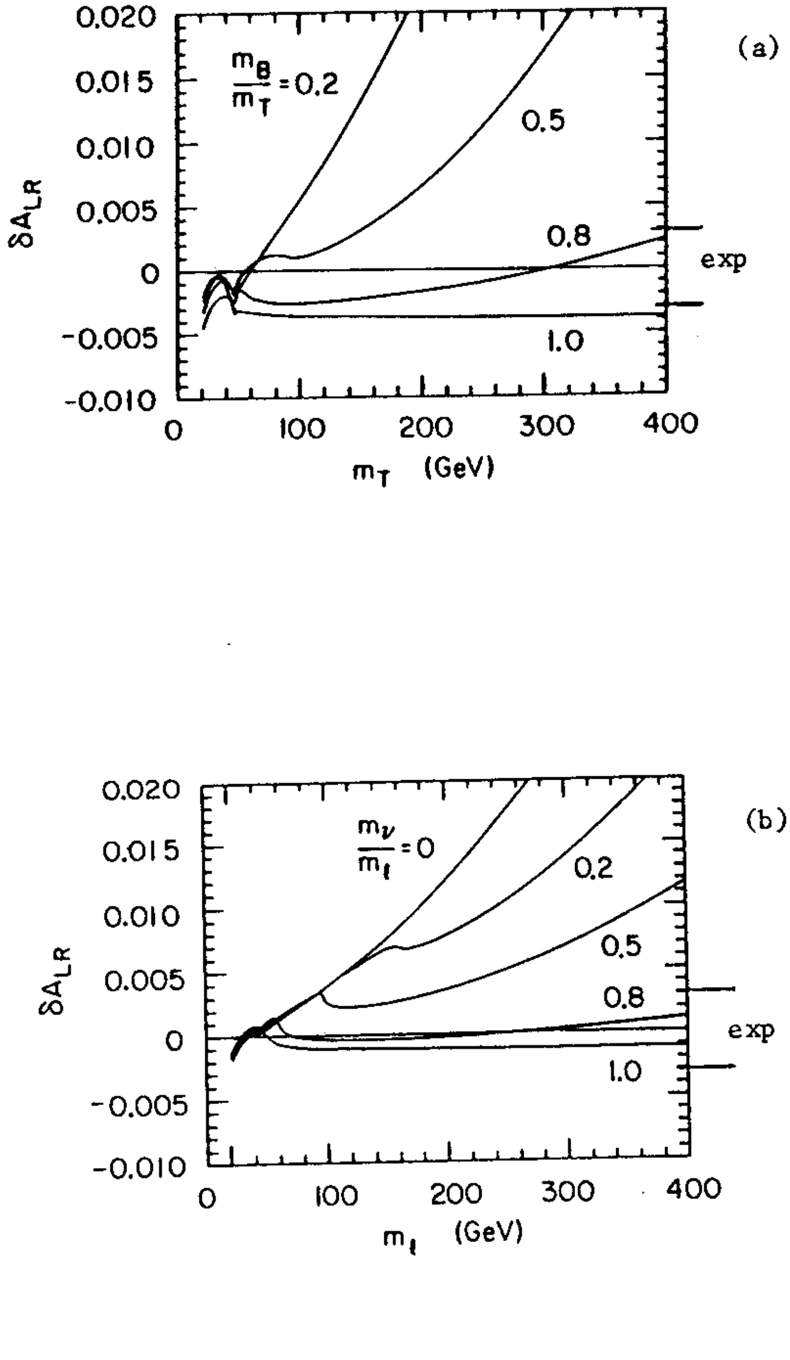


Figure 7: Predicted δA_{LR} for a fourth quark doublet (a) and for a fourth lepton doublet (b) (from [6])

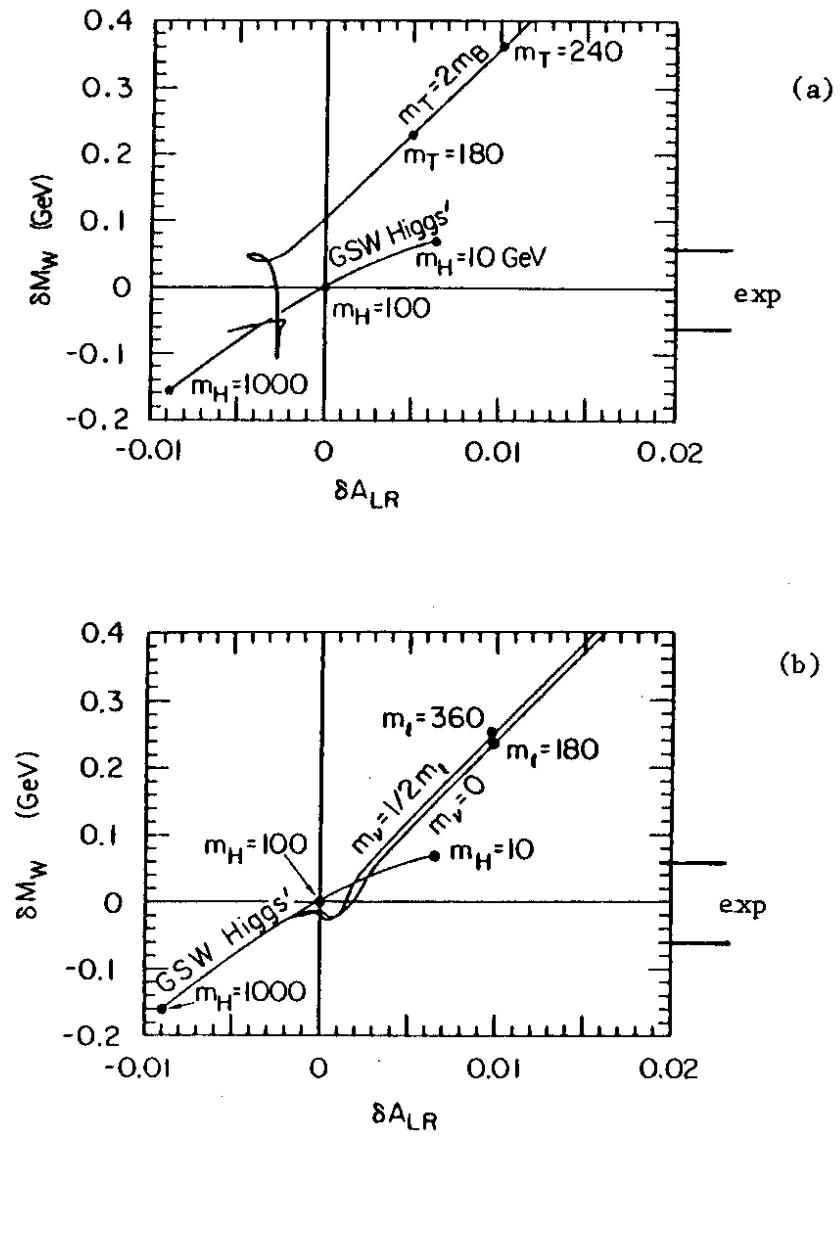
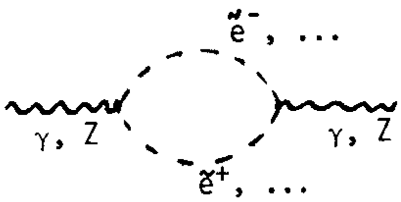


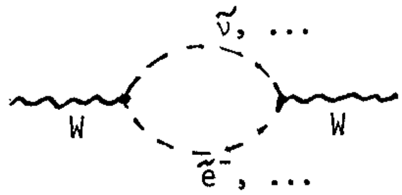
Figure 8: δA_{LR} and M_W for the standard model and for a fourth generation. (a) quark doublet, (b) lepton doublet (from [6])

4.3 The supersymmetric standard model

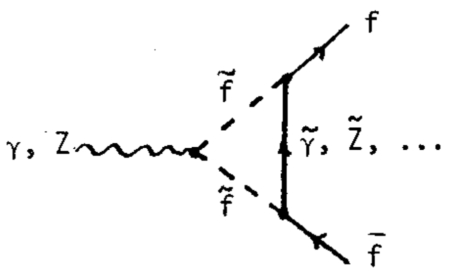
The supersymmetric standard model [22] which in most scenarios of superstring inspired phenomenology arises as the low energy effective theory requires the existence of new superpartners for all known quarks, leptons, gauge and Higgs bosons. Although the structure is to a large extent determined by gauge and supersymmetry there are still quite a few additional free parameters: the masses of the scalar quarks and leptons, and the mixing parameters which diagonalize the gaugino mass matrices.

The universal part of the supersymmetric radiative corrections, that part which contributes exclusively to the vacuum polarizations of the gauge bosons, consists of the squark and slepton loops:





For correctness we have to restrict the $\tilde{q}$ and $\tilde{l}$ in the loops to those which are not the scalar partners of the external fermions in the 4-fermion process under consideration (e.g. $\tilde{l} \neq \tilde{e}$, $\tilde{q} \neq \tilde{f}$ for $e^+e^- \to f\bar{f}$). Otherwise also the non-universal direct corrections like for example



have to be taken into account.

This is also true for the virtual gaugino states (resp. gaugino - Higgsino mixing mass eigenstates), which have direct and hence non-universal couplings to the external fermions. Without referring to specific models the sfermion masses as well as the full mixing structure in the gaugino-Higgsino sector enter the radiative corrections and make a quantitative numerical discussion blown up and involved.

A general discussion of the one-loop radiative corrections to the weak boson masses has been performed in [23] yielding an upper bound of a few hundred MeV in the mass shifts. However, the renormalization scheme applied in [23] is different from the on-shell renormalization followed by most of the actual discussions: the authors of [23] utilize the $\nu_\mu e \to \nu_\mu e$ process for the definition of the mixing angle. The effects in the boson masses hence require the precise knowledge of $\sin^2\theta_W$ derived from the $\nu e/\bar{\nu}e$ cross section ratio which will at best be measured with an accuracy of $\Delta\sin^2\theta_W = \pm 0.005$ [24] and is less precise than from the on-resonance asymmetries.

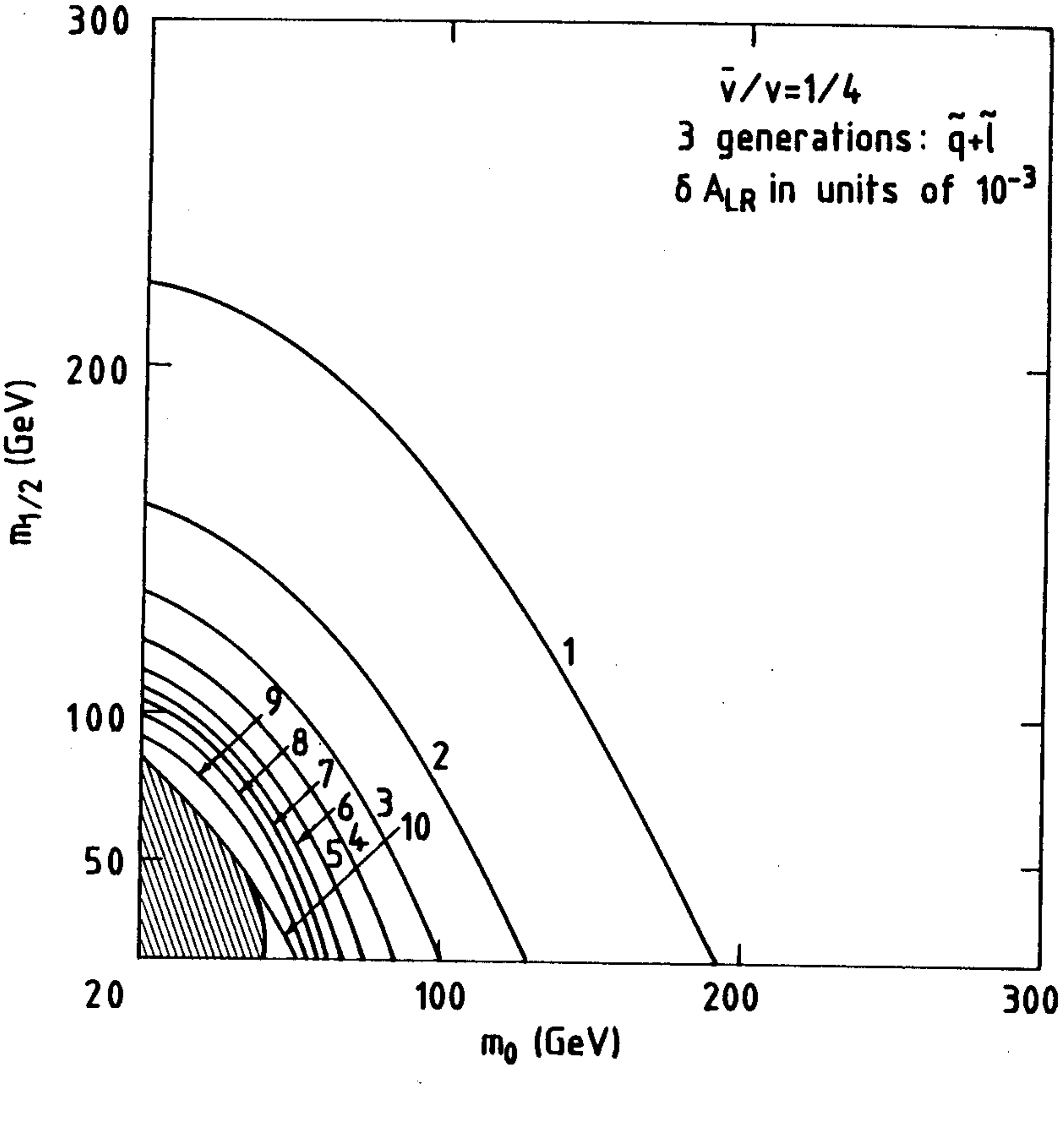


Figure 9: Contour lines of A_{LR} for SUSY sfermions (from [26])

For some specific samples of SUSY particles the on-shell radiative corrections to the left-right asymmetry have been presented in [6]: sfermions with large isospin mass splittings contribute typically of order $\Delta A_{LR} \approx 0.01$, wino states ≈ 0.005.

Sparticle spectra found in realistic models involve also mass splittings between left and right handed sfermions [25]. Representative radiative corrections arising from the vacuum polarizations due to sfermions in phenomenological models have been presented in [26]. All these SUSY corrections turn out to be quite small; they are not visisble in the forward-backward asymmetry but can give rise to observable effects in A_{LR}.

Contour lines for fixed shifts in A_{LR} (in units of 10^{-3}) are shown in Figure 9. They are due to squarks and leptons in the $(m_0,\ m_{1/2})$ plane with the soft SUSY breaking parameters m_0, $m_{1/2}$ of the minimal supersymmetric standard model [25]. Sfermions with masses $\sim$ 100 GeV, accessible in the LEP 200 range, would give a signal comparable to the experimental limit of measurability of $3 \cdot 10^{-3}$. A systematic treatment of the complete SUSY corrections would induce further model dependence via the gaugino sector and has not yet been performed.

4.4 Models with 2 Higgs doublets

The supersymmetric standard model requires the extension of the Higgs sector by a second Higgs doublet and therefore the complete calculation of SUSY radiative corrections includes automatically 2-doublet effects. Nevertheless it is worthwhile to give a seperate discussion of the Higgs sector with two doublets for several reasons:

– Without introducing the concept of supersymmetry the two-doublet model is one of the mildest extensions of the minimal model and has attracked interest by the discussion of CP violation [27] and the Peccei-Quinn mechanism to solve the strong CP problem [28]. The looser restrictions to the Yukawa couplings may give rise to phenomenologically appealing consequences like flavor changing Z decays which are suppressed in the minimal model [29].

– The supersymmetric Higgs sector has very tight constraints to the masses of the scalar bosons which forbid large mass splittings between the $I_3 = +1/2$ and $-1/2$ states. As a consequence, their effects in the radiative corrections are practically invisible. For the non-supersymmetric 2-Higgs doublet model, on the other hand, large radiative corrections can be present [30-31].

The vacuum expectation values v_1, v_2 of the complex doublets ($j = 1, 2$)

$$\Phi_j = \begin{pmatrix} \phi_j^+(x) \\ \phi_j^0(x) \end{pmatrix} = \begin{pmatrix} \phi_j^+(x) \\ (v_j + \eta_j(x) + i\,\chi_j(x))/\sqrt{2} \end{pmatrix} \tag{26}$$

induce the masses of the vectorbosons in the following way:

$$M_W = \frac{1}{2} g_2 \sqrt{v_1^2 + v_2^2}, \quad M_Z = \frac{1}{2}\sqrt{g_1^2 + g_2^2}\,\sqrt{v_1^2 + v_2^2}\,. \tag{27}$$

Three of the eight degrees of freedoms of the doublet fields are absorbed in forming the longitudinal polarization states of the $W^\pm$, Z, and 5 remain as physical particles:

a pair of charged Higgs bosons $H^{\pm}$, two neutral scalars H_0, H_1, and a single neutral pseudoscalar H_2. These physical states are obtained by diagonalizing the mass matrix coming from the Higgs potential:

$$H^+ = -\phi_1^+ \sin\beta + \phi_2^+ \cos\beta \tag{28}$$

$$H_2 = -\chi_1 \sin\beta + \chi_2 \cos\beta$$

for the charged Higgs and the neutral pseudoscalar, and

$$H_0 = \eta_1 \cos\alpha + \eta_2 \sin\alpha \tag{29}$$

$$H_1 = -\eta_1 \sin\alpha + \eta_2 \cos\alpha$$

for the 2 neutral scalars. The mixing angle β is determined by

$$\tan\beta = v_2/v_1 \tag{30}$$

whereas α depends on all parameters of the Higgs potential.

In a non-SUSY 2-Higgs model the angles α,β and all the physical masses M_0,M_1, M_2, M_{H^+} are independent parameters. In the minimal supersymmetric model these quantities are severely constrained [32] :

$$M_{H^+}^2 = M_W^2 + M_2^2$$

$$M_{0,1}^2 = \frac{1}{2}\left(M_Z^2 + M_2^2 \pm \sqrt{(M_Z^2 + M_2^2)^2 - 4\, M_Z^2 M_2^2 \cos^2 2\beta}\right) \tag{31}$$

$$\tan(2\alpha) = \tan(2\beta)\, \frac{M_2^2 + M_Z^2}{M_2^2 - M_Z^2}$$

In such a model one of the neutral scalars is always lighter than the Z, whereas $M_{H^+} > M_W$. From present e^+e^- experiments an experimental lower bound $M_{H^+} > 18$ GeV was derived [33].

If the masses of the Higgs bosons are of the weak boson mass scale or heavier there is little chance to produce them directly in the e^+e^- colliders of the next future. Indirect effects, however, may be present in the radiative corrections to the M_W-M_Z correlation and in $e^+e^- \to f\bar{f}$ around the Z resonance.

4.4.1 The vector boson masses

For the calculation of radiative corrections to fermionic processes we need as additional input parameters only the Higgs masses and the mixing angles α, β.

Application to the μ lifetime yields the relation between M_W and M_Z analogous to (10) :

$$M_W^2(1 - M_W^2/M_Z^2) = \frac{\pi\alpha}{\sqrt{2}G_\mu}\, \frac{1}{1 - (\Delta r_{min} + \Delta\bar{r})} \tag{32}$$

where the minimal radiative correction Δr_{min} is augmented by the non-standard term

$$\Delta\bar{r}(\alpha, M_W, M_Z, M_0, M_1, M_2, M_{H^+}; \mid \alpha - \beta \mid)\,.$$

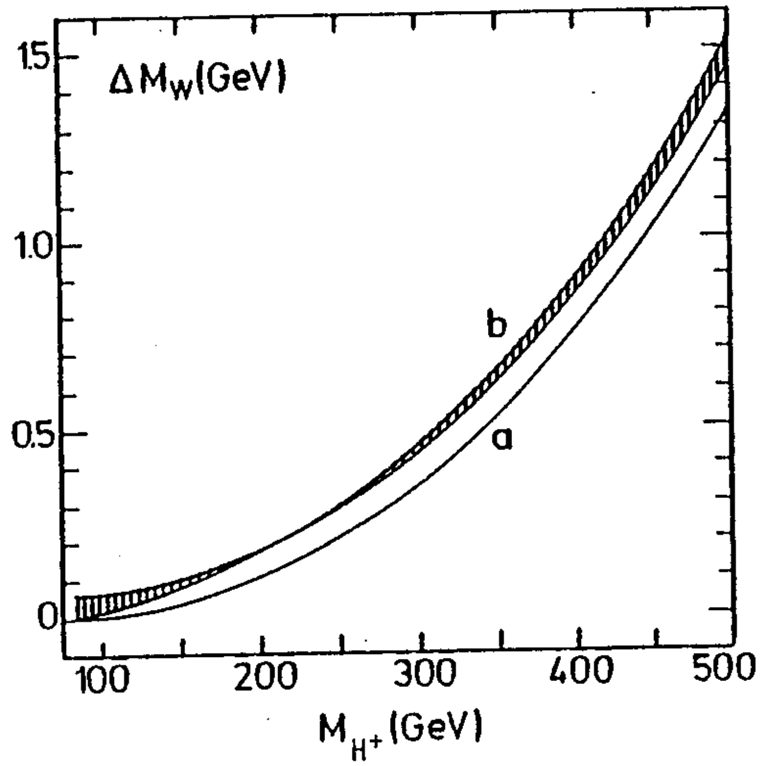


Figure 10:
Additional Higgs contributions to M_W.
$M_Z = 93$ GeV, $M_0 = M_Z$.
a) $M_1 = M_2 = M_Z$
b) $M_1 = 10$ GeV, $M_2 = M_Z$

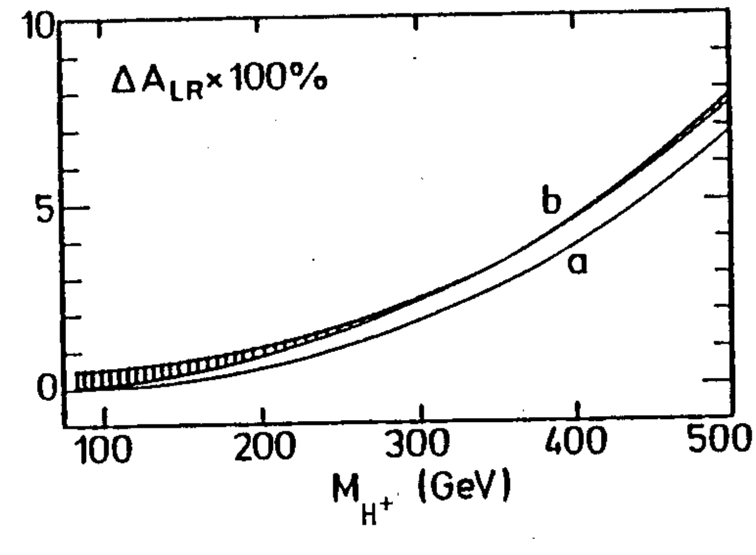


Figure 11a:
Additional Higgs Contributions to
A_{LR} on-resonance
Same parameters as in Figure 10.

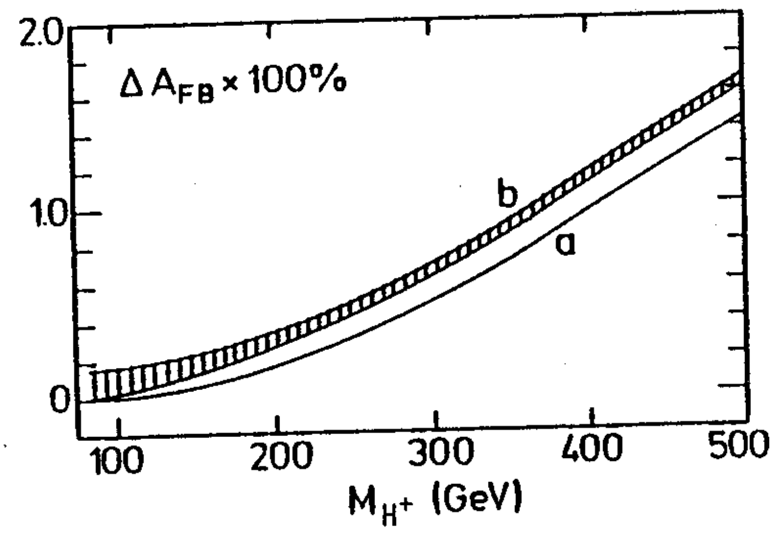


Figure 11b:
Additional Higgs Contributions to
A_{FB} on-resonance
Same parameters as in Figure 10.

$\Delta\bar{r}$ depends on the two mixing angles only in the combination ζ if v_1 and v_2 do not differ by several orders of magnitude.

Here we have to deviate slightly from the strategy summarized in eq. (24): For general mixing in the neutral scalar sector neither H_0 nor H_1 can be identified with the "standard" Higgs. Thus, if one of them (say H_0) is included in the standard Δr it has to be subtracted in $\Delta\bar{r}$.

The numerical solution of (32) for the standard situation with $\Delta r(\alpha, M_W, M_Z, M_H = M_0)$, and for the 2-doublet case with

$$\Delta r_{min} + \Delta\bar{r}(\alpha, M_W, M_Z, M_0, M_1, M_2, M_{H^+}, \zeta)$$

yields the value for M_W after specifying M_Z and the Higgs mass(es) and $\zeta = \mid \alpha - \beta \mid$. The differences are depicted in Figure 10 as functions of the charged Higgs mass (which is assumed to be larger than M_W) for various sets of the neutral scalar/pseudoscalar masses $M_{1,2}$. The shaded area corresponds to the variation of ζ between 0 and $\pi/2$. In the case *a* the result is independent of ζ.

From present M_W measurements with $\Delta M_W = 1.5$ GeV no restrictive bound on the mass splitting between the neutral and charged scalar sector can be derived. An accuracy of $\Delta M_W = 100$ MeV, as expected from LEP 200, can restrict $M_{H^+} < 200$ GeV if the neutral Higgs masses are $< M_Z$.

In the supersymmetric Higgs model the constraints (31) forbid large neutral-charged mass splittings. As a consequence, the deviations from the minimal model remain smaller than the experimental uncertainty in M_W.

4.4.2 Asymmetries

We calculate the $O(\alpha)$ corrected on-resonance asymmetries A_{LR} and $A_{FB}(e^+e^- \to \mu^+\mu^-)$ in the following way:

For given ζ, M_Z and Higgs masses the corresponding value for M_W resp. $\sin^2\theta_W$ is derived from (32). This value is used as input for the calculation of A_{FB} and A_{LR}. Then the standard model result (with $M_H = M_0$) is subtracted yielding the non-standard contributions. These are displayed in Figure 11 for the same set of parameters as in Figure 10. The shaded area indicates the variation with the mixing angle ζ in the neutral Higgs sector.

The left-right asymmetry shows the best sensitivity to mass splittings. An experimental accuracy of $\Delta A_{LR} = \pm 0.003$ can restrict the mass of the charged Higgs boson to $M_{H^+} < 160$ GeV if M_1, M_2 are of the order of the Z mass. The unpolarized forward-backward asymmetry is somewhat less restrictive: from an experimental $\Delta A_{FB} = \pm.002$ a constraint on the charged Higgs mass of $M_{H^+} < 200$ GeV can be obtained. This bound is comparable with that from a W mass measurement with a precision of $\Delta M_W = \pm 100$ MeV (see Figure 10).

Again, for the SUSY Higgs model with the restrictions (31), the absence of large mass splittings keeps the deviations from the minimal model below the experimental sensitivity.

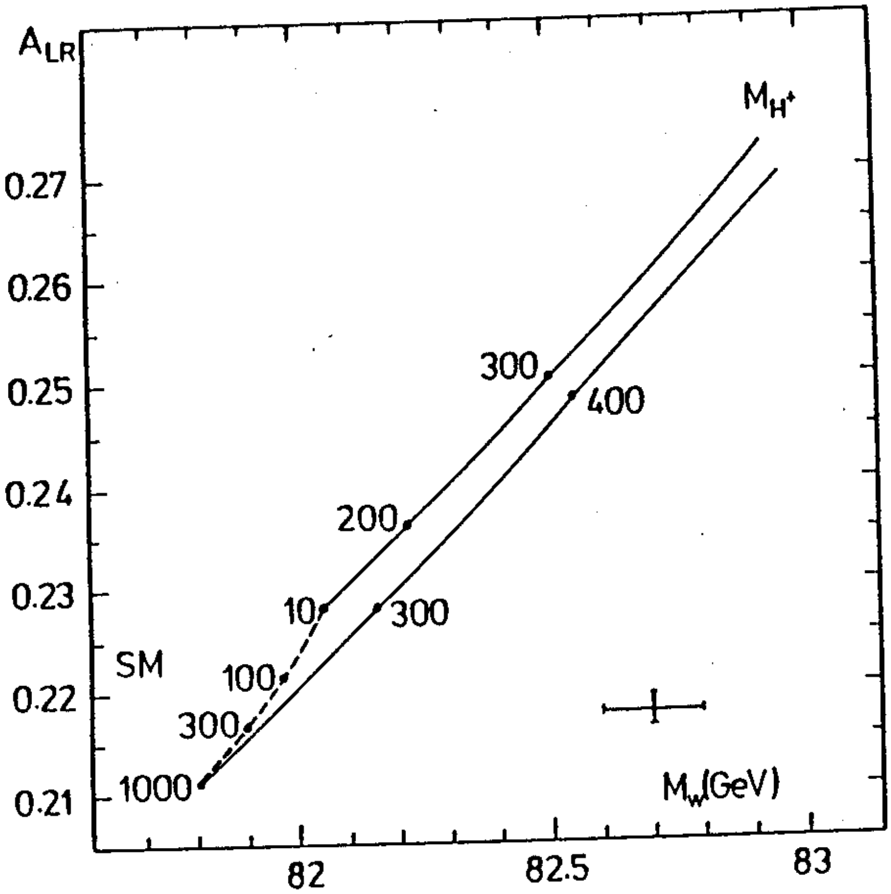


Figure 12: A_{LR} and M_W in the standard model for various M_H (- - - - -) and for an extra Higgs doublet (with $M_0 = M_H$, $M_1 = M_2 = M_Z$).
$M_Z = 93$ GeV, $m_t = 60$ GeV

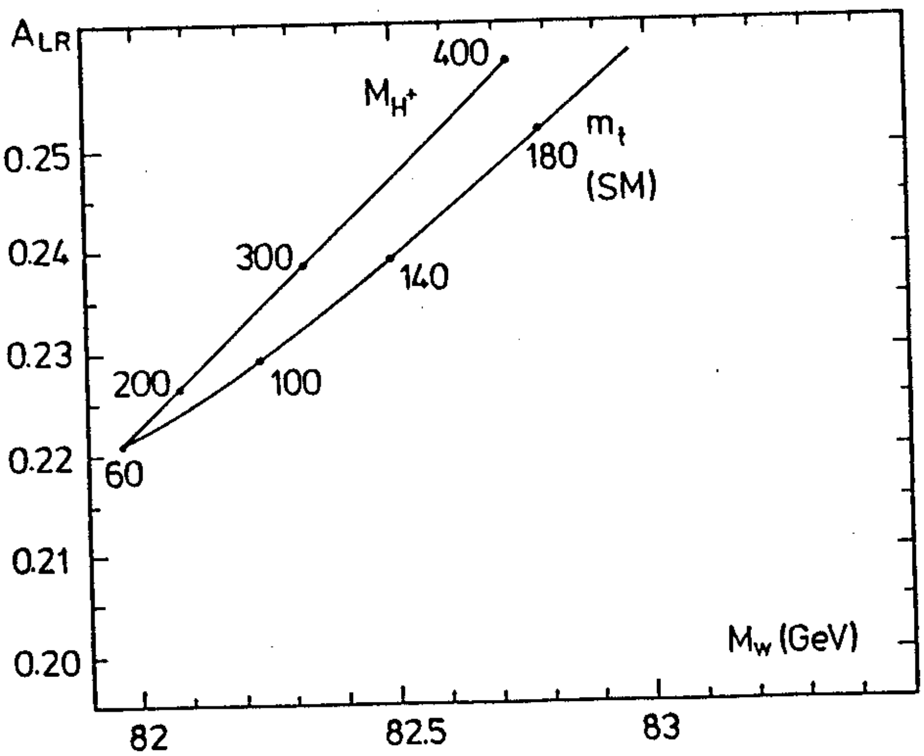


Figure 13: A_{LR} and M_W in the standard model for various m_t and for an extra Higgs doublet (with $M_1 = M_2 = M_Z$, $m_t = 60$ GeV). $M_Z = 93$GeV

A question of particular interest is again the separability from the minimal model. To this end we have plotted in Figure 12 M_W against A_{LR}, for the standard model varying the Higgs mass, and for 2-doublet models where the mass of M_{H^+} is varied and one of the neutral scalars is identified with corresponding standard Higgs. In particular for the situation that a light scalar boson (below 100 GeV) would be discovered the extended model prediction lies in a disjoint region in the parameter plane.

The situation from the top mass uncertainty is less optimistic (Figure 13): The m_t dependence of the minimal model and the dependence on isospin mass splitting in the 2-Higgs model (keeping $m_t = 60$ GeV fixed) are very similar to each other. The additional possibility of a heavy top in the 2-doublet case can bring the two curves quite close to each other. A more restrictive bound on m_t than presently available would therefore become necessary.

4.4.3 Yukawa couplings

In our discussion we have neglected so far the aspects associated with the Yukawa couplings since they are usually small for the known fermions. Models with two Higgs doublets with different vacuum expectation values may develop Yukawa couplings which for the extra neutral and charged Higgs bosons are enhanced by a factor v_1/v_2 resp. v_2/v_1 compared to the standard Yukawa couplings.

Leptons:

If we assume that the scalar doublet Φ_1 is responsible for the masses of the charged leptons the corresponding Yukawa term with coupling constant g_l reads for one family:

$$L_{lept}^{Yu} = -g_l\,(\bar{\nu}_L\,\phi_1^+\,l_R + \bar{l}_R\,\phi_1^-\,\nu_L + \bar{l}_L\,\phi_1^0\,l_R + \bar{l}_R\,\phi_1^{0*}\,l_L)\,. \tag{33}$$

Identifying the lepton mass by $m_l = g_l\,v_1/\sqrt{2}$ the couplings of the physical Higgs states follow from (33) after the rotation (28,29):

$$\begin{aligned} L_{lept}^{Yu} \;=\; & \frac{g_2}{\sqrt{2}}\cdot\frac{m_l}{M_W}\cdot\tan\beta\left(\bar{\nu}\frac{1+\gamma_5}{2}l\cdot H^+ \;+\; \bar{l}\frac{1-\gamma_5}{2}\nu\cdot H^-\right) \\ & + \frac{g_2}{2}\cdot\frac{m_l}{M_W}\cdot\tan\beta\cdot\bar{l}\,i\gamma_5\,l\cdot H_2 \\ & - \frac{g_2}{2}\cdot\frac{m_l}{M_W}\cdot\bar{l}\,l\left(\frac{\cos\alpha}{\cos\beta}H_0 - \frac{\sin\alpha}{\cos\beta}H_1\right)\,. \end{aligned} \tag{34}$$

For $v_1 \approx v_2$, i.e. for $\tan\beta = O(1)$, the coupling strength is similar as in the minimal model. The special situation $v_1 \ll v_2$, where the boson masses are essentially determined by the vacuum expectation value of Φ_2 only, leads to an enhancement of the H^+ and H_2 couplings by the factor $\tan\beta = v_2/v_1$. The second (neutral) mixing angle α behaves like[3]

$$\tan(2\alpha) \sim \frac{v_1}{v_2} \to 0 \quad (\alpha \approx 0 \text{ or } \alpha \approx \frac{\pi}{2})\,.$$

[3]at least if the quartic couplings in the Higgs potential are all of the same order of magnitude

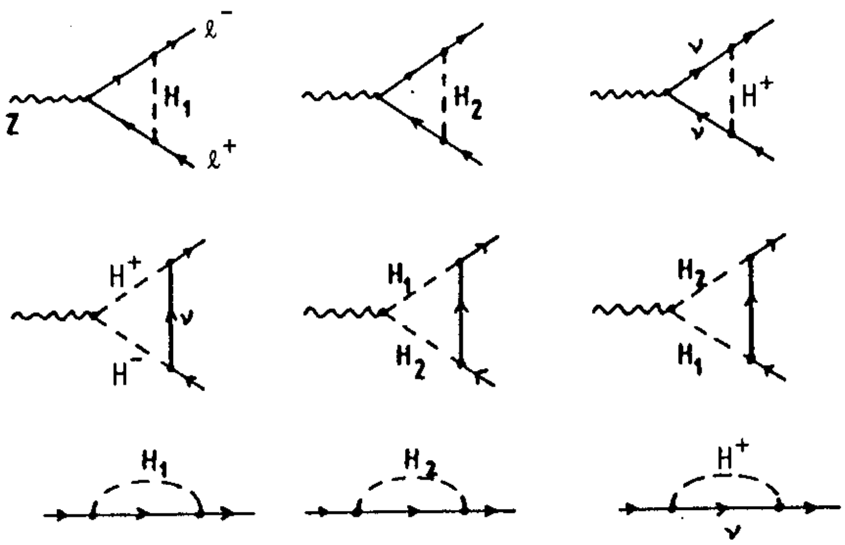


Figure 14: Extra Higgs contributions to the $Z\tau\tau$ vertex

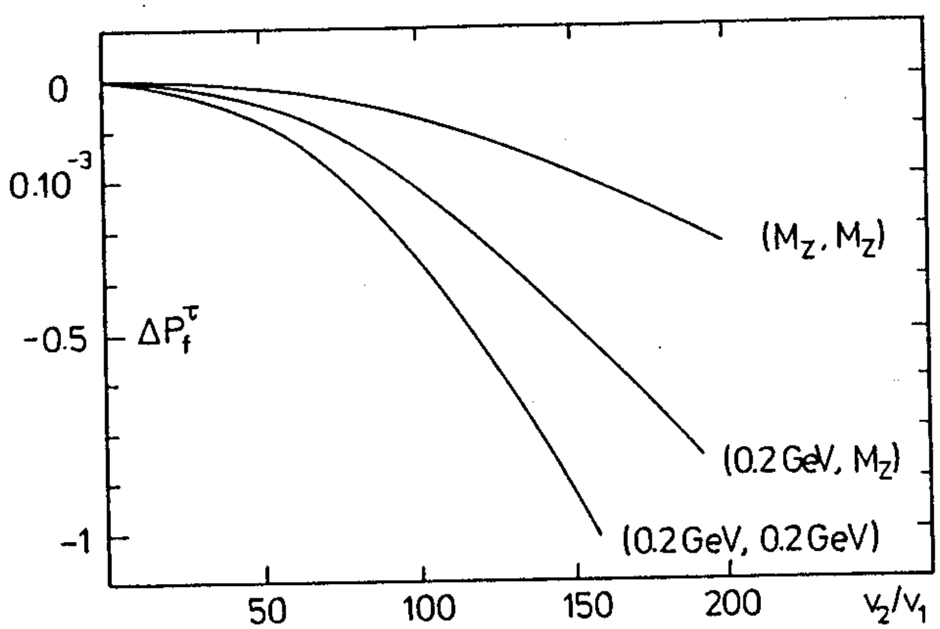


Figure 15: τ final state polarization for enhanced Yukawa couplings. M_1 and M_2 in parantheses, $M_{H^+} = 100$ GeV

If we choose the possibility $\alpha \approx \pi/2$ then $\cos\alpha/\cos\beta$ is of order 1 and hence the H_0 coupling is negligible for the light leptons. H_1, on the other hand, gets an enhancement similar to $H^{\pm}$, H_2.

The appearence of $H^{\pm}$ and $H_{1,2}$ in the vertex corrections to $Z - l - l$ $(\gamma - l - l)$, displayed in Figure 14, yields different contributions for μ and τ final states in $e^+e^- \rightarrow l^+l^-$ and therefore to a small possible deviation from the leptonic universality. Without information from the hadronic part of the Yukawa interaction very little constraint on $\tan\beta$ is experimentally known. The only relevant restriction comes from the magnetic moment $g-2$ of the muon [34] :

$$\frac{v_2}{v_1} < 70 \tag{35}$$

under the assumption that the additional neutral scalar/pseudoscalar pair is not degenerate in mass (the value was derived for masses 6 GeV and 100 GeV [34]). Mass degeneracy leads to significantly weaker bounds than (35).

A measurement of the $Z - \tau$ vertex in $e^+e^- \rightarrow \tau^+\tau^-$ on the Z resonance offers another possible investigation of the additional scalar contributions, either in terms of A_{FB}^{pol} or by a measurement of the final τ polarization. Both quantities are determined by A_τ from (20) which is sensitive to the vertex contributions of Figure 14.:

$$v_\tau \rightarrow v_\tau + \Delta F_V^\tau , \quad a_\tau \rightarrow a_\tau + \Delta F_A^\tau .$$

The effects of the additional Higgs particles on A_τ are shown in Figure 15 for various values of the mass parameters. Precision measurements on τ polarization yields leptonic information on $\tan\beta$ supplementary to $g-2$; for mass degenerate neutral scalar/pseudoscalar states the information is even complementary, although not very restrictive.

Quarks:

The structure of the Yukawa couplings that would arise from a supersymmetric model implies that Φ_1 gives rise to the down quark mass and Φ_2 to the up quark mass. A non-SUSY argument suggesting such a pattern is the absence of flavor changing neutral currents.

The Higgs-quark interaction term is conveniently written in the doublet field components (26) (neglecting quark mixing angles):

$$\begin{aligned} L_{quark}^{Yu} \;=\; & - g_d\,(\overline{u}_L\,\phi_1^+\,d_R + \overline{d}_R\,\phi_1^-\,u_L + \overline{d}_L\,\phi_1^0\,d_R + \overline{d}_R\,\phi_1^{0*}\,d_L) \\ & - g_d\,(\overline{u}_R\,\phi_2^+\,d_L + \overline{d}_L\,\phi_2^-\,u_R + \overline{u}_R\,\phi_2^0\,u_L + \overline{u}_l\,\phi_2^{0*}\,u_R) \end{aligned} \tag{36}$$

Identification of the u and d masses as

$$m_u = \frac{g_u}{\sqrt{2}}v_2, \quad m_d = \frac{g_d}{\sqrt{2}}v_1$$

we find the couplings for the physical states:

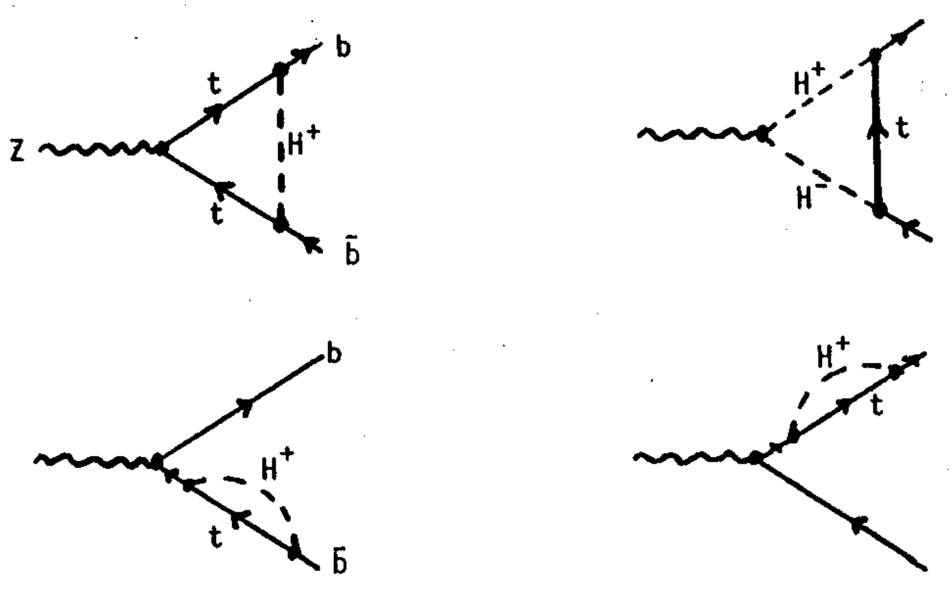
$$
\begin{aligned}
L^{Yu}_{quark} \;=\; & \frac{g_2}{\sqrt{2}} \left(\frac{m_d}{M_W} \tan\beta \cdot \bar{d} \frac{1+\gamma_5}{2} u + \frac{m_u}{M_W} \cot\beta \cdot \bar{u} \frac{1-\gamma_5}{2} d \right) H^+ + h.c. \\
& + \frac{g_2}{2} \left(\frac{m_d}{M_W} \tan\beta \cdot \bar{d}\, i\gamma_5\, d + \frac{m_u}{M_W} \cot\beta \cdot \bar{u}\, i\gamma_5\, u \right) H_2 \\
& - \frac{g_2}{2} \cdot \frac{m_d}{M_W} \cdot \bar{d}\, d \left(\frac{\cos\alpha}{\cos\beta} H_0 - \frac{\sin\alpha}{\cos\beta} H_1 \right) \\
& - \frac{g_2}{2} \cdot \frac{m_u}{M_W} \cdot \bar{u}\, u \left(\frac{\sin\alpha}{\sin\beta} H_0 + \frac{\cos\alpha}{\sin\beta} H_1 \right) . \qquad (37)
\end{aligned}
$$

Two different scenarios yielding a set of enhanced Yukawa couplings in (37) are possible: The situation $v_2 \gg v_1$, in analogy to the leptonic case, which enhances the d-type couplings to $H^{\pm}$, $H_{1,2}$ by v_2/v_1, and the other possibility of having $v_1 \gg v_2$ enhancing the u-type couplings by v_1/v_2.

A situation of particular interest also for the (non enhanced) case $\tan\beta = 1$ ($v_1 = v_2$) is encountered in the $Z - bb$ vertex which can be investigated on the Z peak also with high precision. The charged Higgs coupling to the (t, b) family involves the term

$$
\frac{g_2}{\sqrt{2}} \cdot \frac{m_t}{M_W} \cdot \frac{v_1}{v_2} \cdot \frac{(1 \pm \gamma_5)}{2}
$$

which can become large if the top quark is heavy. The $H^{\pm}$ bosons in connection with virtual t quarks enter the $Z - bb$ vertex as follows:



The partial decay width $\Gamma_Z(Z \to b\bar{b})$ is displayed in Figure 16 for the minimal model and in the presence of a charged Higgs boson (with $M_{H^+} = 100$ GeV). The standard model result is practically independent of m_t [35]. The reason for this behaviour is the top dependence of the vertex corrections which cancels (to a large extent) the top contributions in the gauge boson 2-point functions. The importance of the t quark in the vertex corrections is underlined when the $Z \to d\bar{d}$ width is considered: its increase for large m_t is a consequence of the propagator corrections. A similar increase is expected for a next quark family with $t' - b'$ mass splitting.

The top independence of the $Z \to b\bar{b}$ width in the minimal model recommends this quantity as a sensitive probe for new particles with direct couplings to the top quark like charged Higgs bosons.

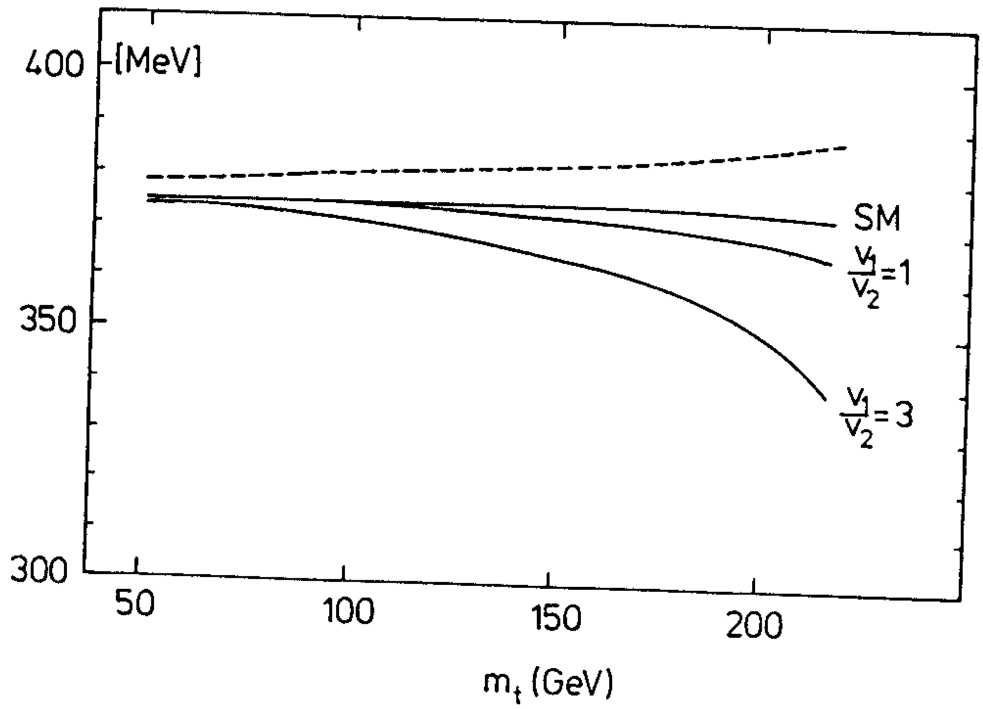
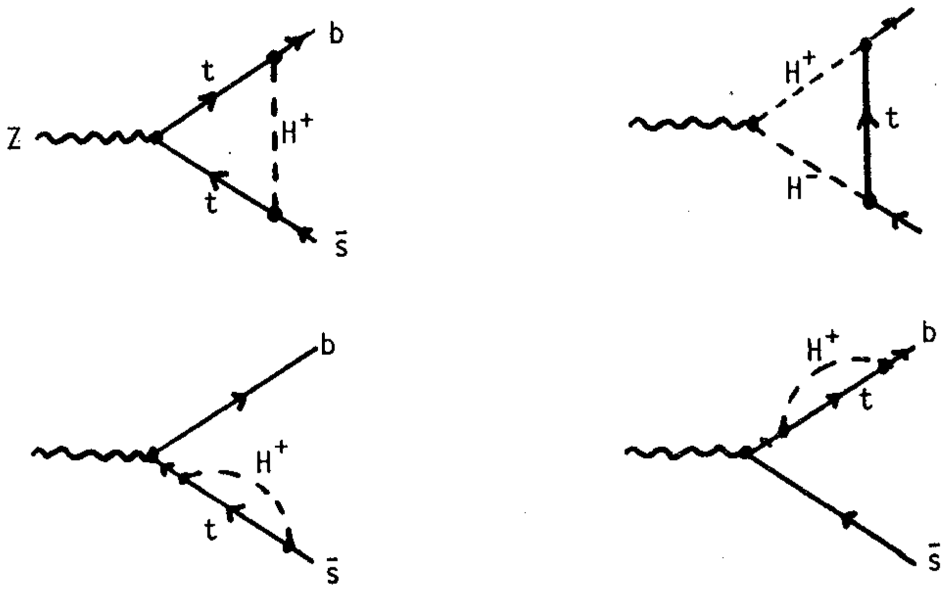


Figure 16:
$\Gamma_Z(Z \to d\bar{d})$ (- - - - -) in the standard model and $\Gamma_Z(Z \to b\bar{b})$ (——) in the standard model (SM) and in the presence of charged Higgs bosons.

Flavor changing Z decays:

The presence of charged Higgs bosons and the possibility $v_1 > v_2$ may also enhance the flavor changing Z decay rates like $Z \to b\bar{s}$ which are too small to be observed in the minimal model [29]. The following diagrams with $H^\pm$ contribute to $Z \to b\bar{s}$, where only the top contribution is significant:



The branching ratio

$$\frac{\Gamma(b\bar{s} + \bar{b}s)}{\Gamma_Z} = B$$

depends on the factor

$$\mid U_{ts} \mid^2 \cdot \left(\frac{m_t}{M_W} \frac{v_1}{v_2} \right)^4$$

A recent discussion in [36] shows that values for v_1/v_2 between 2 and 5 yield branching ratios $B > 10^{-6}$ which may become experimentally observable.

5 New particles in SU(2)×U(1) not preserving the tree level structure in 4-fermion processes

5.1 General Higgs structure

All kinds of new physics encountered so far share the following common features with the standard model:

- $\sin^2 \theta_W$, which measures the admixture of the electromagnetic current to the weak neutral current is determined when M_Z and M_W are known.

- $\sin^2\theta_W$, if measured from ν scattering or from the on-resonance asymmetries, is equal to $1 - (M_W^2/M_Z^2)_{exp}$ when radiative corrections are taken into account properly.

The tree level structure of the interaction between gauge bosons and fermions is given by the interaction Lagrangian

$$L = \frac{g_2}{\sqrt{2}}(J_+W^- + J_-W^+) + \frac{g_2}{\cos\theta_W} J_N Z^0 + (g_2 \sin\theta_W) J_{em} A \tag{38}$$

where the currents are built from the fermion doublets, isospin matrices and $Q = I_3 + Y/2$ as

$$\begin{aligned} J_\pm^\mu &= \bar{\psi} I_\pm \gamma^\mu \frac{1-\gamma_5}{2} \psi , \\ J_N^\mu &= \bar{\psi} \left(I_3 \gamma^\mu \frac{1-\gamma_5}{2} - \sin^2\theta_W Q \gamma^\mu \right) \psi , \end{aligned} \tag{39}$$

and the electric charge is identified by $e = g_2 \sin\theta_W$.

In general the mixing angle diagonalizes the neutral boson mass matrix and is determined by the vacuum structure of the Higgs sector (if we exclude other scalar particles like sfermions to involve non-zero vacuum expectation values). If only Higgs doublets (and eventually singlets) have non-vanishing vacuum expectation values the consequence $\cos\theta_W = M_W/M_Z$ allows to predict M_W from the parameters of the neutral current sector.

In the general case, however, this is no longer possible, but the structure of the interaction Lagrangian (38) ist still valid. We only have to give up the tight connection between the vector boson masses and the mixing angle and allow for an additional independent parameter which connects the neutral and charged current sector. Conveniently this parameter is chosen to be

$$\rho = \frac{M_W^2}{M_Z^2 \cos^2\theta_W} . \tag{40}$$

It can be defined e.g. in terms of the NC/CC ration in νN scattering with isoscalar targets. Im view of the high precision experiments associated with asymmetry measurements on the Z peak it may be useful to treat $\sin^2\theta_W$ as a quantity independent of M_W, M_Z

$$\sin^2\theta_W = 1 - \frac{M_W^2}{\rho M_Z^2}$$

measurable e.g. from A_{LR} by means of (20). The structure of the tree level coupling constants in (1) resp. (20) is preserved if expressed by means of $\sin^2\theta_W$ since its genuine meaning in (38,39) is not changed. All other tree level formulae can easily be adapted by performing the transformation

$$M_Z \to \sqrt{\rho} M_Z .$$

For a general Higgs sector of the SU(2)×U(1) theory the deviation from $\rho = 1$ can be written in terms of the Higgs representations ϕ_i which have non-zero vacuum expectation values $< \phi_i >$:

$$\rho = 1 + \frac{\sum_i < \phi_i \vec{I}^2 \phi_i > - 3 \sum_i < \phi_i I_3^2 \phi_i >}{2 \sum_i < \phi_i I_3^2 \phi_i >} . \tag{41}$$

For a discussion at the tree level the information collected in ρ is sufficient. For a discussion at the level of radiative corrections more detailed information on the non-doublet Higgs representations would become necessary, e.g. the masses of the new Higgs particles entering the gauge boson propagators.

The general expression of a measurable quantity

$$A = A_{tree} + \Delta A(\rho = 1) + \Delta A(\rho \neq 1) \tag{42}$$

involves radiative corrections which can be split up in the following way:

The part $\Delta A(\rho = 1)$ is identical with that of the minimal model (or with extensions not changing $\rho = 1$). The additional term $\Delta A(\rho \neq 1)$ contains the virtual contributions of all new objects which cause a deviation from $\rho = 1$. This last term has not yet been included in loop calculations for the general case $\rho \neq 1$. Usually it is assumed that that it can be neglected, i.e. that the tree level effects of the new structure are bigger than the induced loop effects. This assumption need not be correct; it is supported, however, by the consistency check that the experimental value for ρ obtained in this way is close to unity. There is presently no obvious experimental indication for $\rho \neq 1$ [37] :

$\rho = 1.002 \pm 0.012 \pm 0.007$ from νN scattering,

$\rho = 1.02 \pm 0.05$ from PETRA/PEP data.

If we rely on the assumption that we can treat the radiative corrections as in the minimal model we get a handle on ρ by the two sets of independent experiments:

1. M_W, M_Z, $\sin^2\theta_W$:
   The direct measurements of the Z and W mass and of $\sin^2\theta_W$ from A_{LR}

   $$A_{LR} = \frac{2(1 - 4\sin^2\theta_W)}{1 + (1 - 4\sin^2\theta_W)^2} + \Delta A_{LR}(\rho = 1)$$

   with the standard model radiative correction $\Delta A_{LR}(\rho = 1)$ allows to determine deviations from 1 within $\Delta\rho < 0.0026$. Thereby it is assumed that the following experimental accuracy can be obtained:
   $\Delta M_Z = 20$ MeV, $\Delta M_W = 100$ MeV, $\Delta \sin^2\theta_W = 0.0004$.

2. M_z, G_μ, $\sin^2\theta_W$:
   The modification of the $M_W \leftrightarrow M_Z$ correlation (10)

   $$M_W^2 \left(1 - \frac{M_W^2}{\rho M_Z^2}\right) = \frac{A}{1 - \Delta r} , \quad A = \frac{\pi\alpha}{\sqrt{2} G_\mu} \tag{43}$$

allows to derive ρ from the measured values $\sin^2\theta_W$, M_Z together with G_μ:

$$\rho = \frac{A}{M_Z^2 \sin^2\theta_W \cos^2\theta_W (1-\Delta r)} .$$

Assuming the same accuracy for the mass measurements as in 1. we expect to be sensitive to $\rho \neq 1$ within $\Delta\rho < 0.002$.

In both cases the experiments are necessary to define the input parameters if we allow for a more general scalar sector. At the same time a deviation $\Delta\rho$ from 1 represents a global measure of a new structure without specific model assumptions.

A test of the hypothesis that the underlying model is of the SU(2)×U(1) type, fixed by independent quantitites α, M_W, M_Z, ρ, can be performed by a combination of 1 and 2: the value for ρ derived from 2 has to reproduce the value for $\sin^2\theta_W$ as determined according to 1, and vice versa. This will become of importance if a significant deviation from the canonical value $\rho = 1$ would be observed. Further strategies to delimit new structures within SU(2)×U(1) from new currents extending SU(2)×U(1) will be discussed in section 6.

5.2 Excited vector bosons

The previous discussion of possible new structures was aligned with the architecture of a SU(2)×U(1) gauge field theory and electroweak symmetry breaking via the Higgs mechanism, allowing a perturbative treatment around the tree level phenomenology. On the other hand, a large Higgs mass $M_H > 1$ TeV corresponds to a strongly interacting bosonic sector within SU(2)×U(1) where perturbation theory becomes unreliable [38]. Although the conventional perturbative influence of M_H in 4-fermion processes remains small also for a very heavy Higgs [43], the appearance of new vector boson bound states which mix with W and Z gives rise to distortions of the phenomenology in the LEP energy domain.

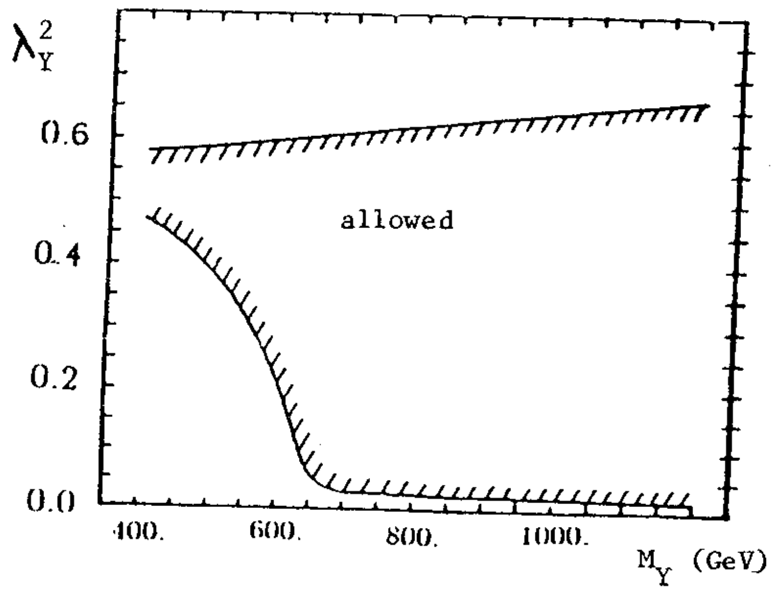


Figure 17: Allowed mass and coupling range for a Y boson resulting from A_{LR} and A_{FB} at $\sqrt{s} = M_Z$ (from [39])

New vector boson states coupling to the weak hypercharge current are usually contained also in models where W and Z are composite objects. A discussion of the influence of those isoscalar vector bosons on the on-resonance asymmetries is given in [39] using an effective tree level Lagrangian approach. Boundaries limiting the range of the isoscalar boson mass M_Y and coupling $g_Y = e/\lambda_Y$ as they would follow from a measurement of A_{LR} and $A_{FB}(e^+e^- \to \mu^+\mu^-)$ are illustrated in Figure 17. They correspond to experimental errors of 1% resp. 2% [39]. A similar analysis concerning new vector states from a strongly interacting Higgs sector has been performed in [40].

6 New currents and extra gauge bosons

Many extensions of the standard model predict the existence of additional currents and new gauge bosons associated with them. Augmentations of the gauge group have been studied for various motivations:

– A possible left - right symmetric structure based on $SU(2)_L \times SU(2)_R \times U(1)$, where parity invariance is broken spontaneously [17]. It involves right handed $W^\pm$ and an additional neutral vector boson.

– A previous E_6 symmetry as in many models inspired by the superstring theory. It allows as the minimal possibility the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{Y'}$ as the effective low energy symmetry [18]. The extra neutral gauge boson coupling to the extra hypercharge Y' can mix with the "old" Z^0 to form physical mass eigenstates Z, Z' where the mass of the lower lying state is in the same range as the standard Z mass.

Since on top of this resonance the direct effect of the Z' with the higher mass is suppressed by the Z' propagator deviations from the standard model come mainly from the admixture of the new current in the physical state Z which modifies the standard neutral current according to

$$J_{N_1} \sim J_{3L} - \sin^2\theta_W\, J_{em} + \frac{g_1'}{g_1}\tan\theta\, J_{Y'} \,. \tag{44}$$

In the on-resonance asymmetries the normalization of J_{N_1} drops out, but the $\sin^2\theta_W$ dependence is distorted by the Y' current. The admixture in (44) depends on the details of the specific model: on the ratio g_1'/g_1 of the $U(1)_Y$ and $U(1)_{Y'}$ gauge couplings,the $Z^0 - B'$ mixing angle θ, and the hypercharge quantum numbers Y' of the fermions. Many phenomenological applications adopt the parameters according to E_6 models related to string theory [18-20].

For the simplest case of one extra neutral boson the fermion - gauge boson interaction Lagrangian has the form

$$\begin{aligned} L &= \frac{g_2}{\sqrt{2}}\left(J_+W^- + J_-W^+\right) \\ &\quad + g_2\, J_{3L}W^3 + g_1\, J_Y B + g_1'\, J_{Y'}B' \,. \end{aligned} \tag{45}$$

The charged current part is identical to that of the minimal model in (38), the neutral current sector differs formally by the additional Y' term from the standard model neutral current in (38). The mass eigenstates are mixtures of the Z^0 and B':

$$\begin{pmatrix} Z \\ Z' \end{pmatrix} = \begin{pmatrix} \cos\theta & \sin\theta \\ -\sin\theta & \cos\theta \end{pmatrix} \begin{pmatrix} Z^0 \\ B' \end{pmatrix} \tag{46}$$

and the mixing angle can be expressed in terms of the physical masses:

$$\tan^2\theta = \frac{M_{Z^0}^2 - M_Z^2}{M_{Z'}^2 - M_{Z^0}^2}, \tag{47}$$

Z^0 being the usual $W^3 - B$ linear combination

$$Z^0 = \cos\theta_W\, W^3 + \sin\theta_W\, B$$

and M_{Z^0} related to M_W by means of (43). For a very heavy Z' the mixing vanishes, and the neutral current according to Z becomes identical with that of the minimal model. Conventionally the lower mass state, denoted by Z, will be identified with the particle observed in $p\bar{p}$ collisions, and, in future, in e^+e^- annihilation.

For a rigorous discussion of the presumably small mixing effects induced in the asymmetries the treatment of the complete one-loop corrections in the extended model seems desirable. At present no such discussion does exist, and we will hence restrict ourselves to the same assumption as in section 5: that the loop contributions are dominated by the minimal model radiative corrections and new objects can be neglected in the loop diagrams in comparison with their tree level effects. The question of radiative corrections in the extended gauge group SU(2)×U(1)×G is adressed to some extent in [20] claiming that the neglected effects are all of order

$$\frac{\alpha}{\pi}\left(\frac{M_Z}{M_{Z'}}\right)^2 \to 0$$

if the mass of the second neutral boson is sufficiently large.

6.1 On-resonance asymmetries and new Z bosons

The possible existence of an extended electroweak gauge structure has induced a lot of activities in investigating the response of the precisely measurable asymmetries in $e^+e^- \to f\bar{f}$ to various concrete realizations of $SU(2)\times U(1)_Y \times U(1)_{Y'}$ models [18-20]. Model independent formulae for the $e^+e^- \to f\bar{f}$, e^+e^- cross sections with general polarizations in the presence of several neutral bosons have been given already in [41].

Since a comprehensive discussion of all possible situations of physical interest would be outside the scope of this article we want to restrict ourselves to the "minimal superstring inspired model" [18] as a representative example and the effects induced in the left-right asymmetry as the most promising place to look for. The gauge boson masses are generated by two Higgs doublets (vacuum expectation values v and $\bar{v}$) and one Higgs singlet (with vacuum expectation value x). With this Higgs structure the mixing angle θ and the mass of the heavier Z' boson are the only additional free parameters if M_Z and $\sin^2\theta_W$ have been fixed. The ratio of the Higgs vacuum expectation values are

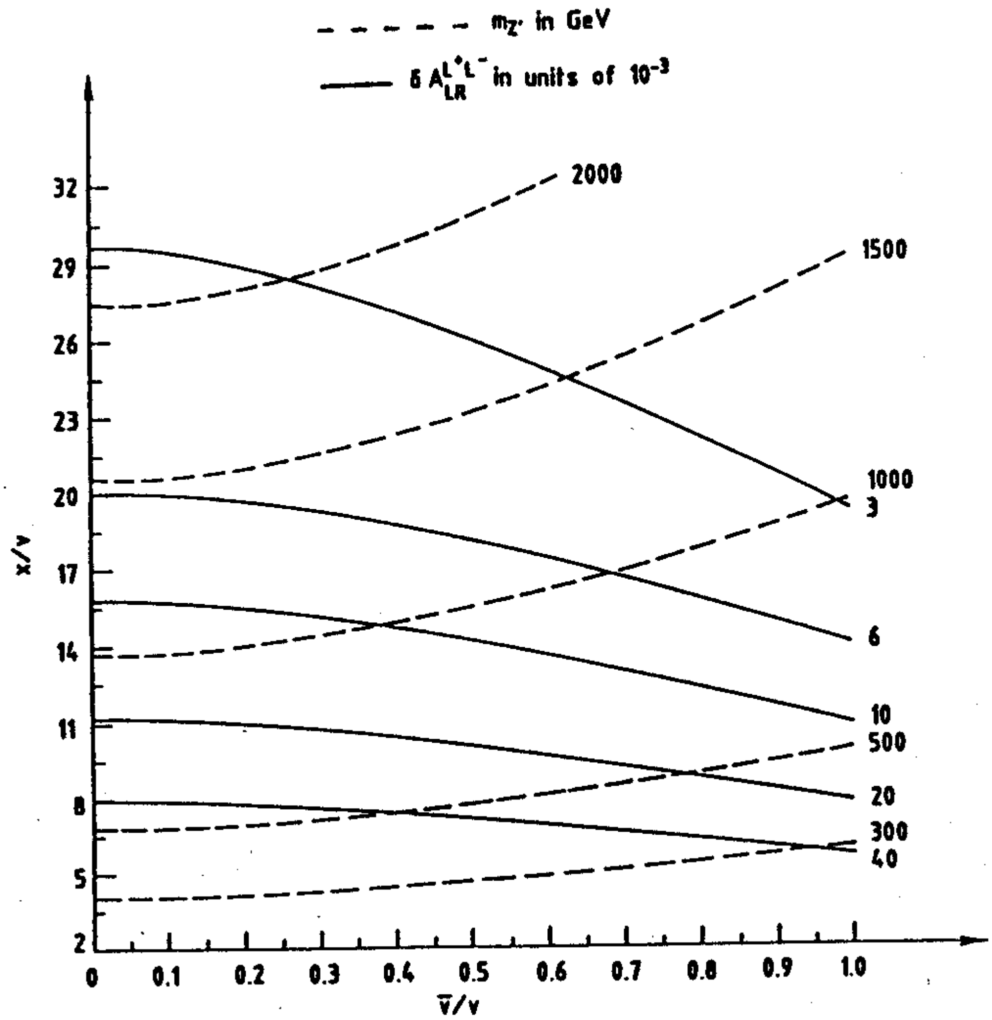


Figure 18: Contour lines for δA_{LR} and $M_{Z'}$ in the minimal superstring inspired model (from [18])

the input on the axis of Figure 18 showing contour lines of constant deviations δA_{LR} (full lines) from the standard model. They visualize a large domain in the parameter space that can be covered experimentally ($\delta A_{LR} < 0.003$), physically more instructively expressed in terms of the second Z' mass (dashed lines). A precision measurement of A_{LR} can therefore be sensitive to a Z' far above the energy range of LEP 200.

6.2 Isolating new Z bosons

In general the effects of different kinds of new physics cannot be identified separately. Deviations from the standard model might be observed, but their origin cannot be isolated unambigously. A combination of various measurable quantities, however, may be helpful to disentangle specific aspects of new contributions beyond the minimal model. In [42] it is discussed how to identify the Z' boson associated with an E_6 unifying group by means of polarized forward-backward asymmetries for various types of fermions.

A question of particular importance is how to distinguish between virtual effects of new objects in SU(2)×U(1) and those effects which are due to a more complicated electroweak gauge group. Although this question cannot be answered in full generality, a strategy for isolating Z' effects has been proposed by Cvetič and Lynn [20]. The method is based on a specific combination of the left-right asymmetry A_{LR} and the polarized forward-backward asymmetry $A_{FB}^{pol}(e^+e^- \to f\bar{f})$ which is insensitive (at the one-loop level) to all new particles in SU(2)×U(1) with couplings to the gauge bosons but not to the external fermions. Those objects would influence the value of the asymmetries exclusively via the vacuum polarization functions (22). Approximately the method can be extended also to particles with direct couplings to the fermions if the additional loop contribution can be neglected compared to the effect of an extra Z' boson. In most cases those direct contributions are smaller (of order $\sim 10^{-3}$) than the experimental error for $A_{FB}^{pol}(e^+e^- \to f\bar{f})$.

The strategy can be easily understood from the following properties of the on-resonance asymmetries (neglecting γ exchange):

* Both A_{LR} and A_{FB}^{pol} are determined by the combination A_f, eq. (20):

$$\begin{aligned} A_{LR} &= A_e \\ A_{FB}^{pol}(e^+e^- \to f\bar{f}) &= \frac{3}{4} A_f \quad (f \neq e)\,. \end{aligned} \tag{48}$$

* The quantity A_f, with any kind of radiative corrections in SU(2)×U(1) of the "oblique" type [4], can be written as

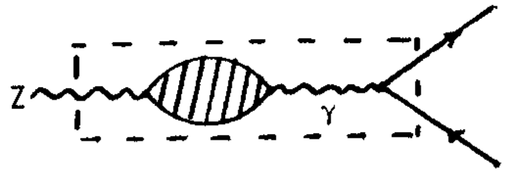
$$Af = \frac{2 I_3^f \, (I_3^f - 2\sin^2\theta_W Q_f - 2Q_f \cdot \Delta)}{(I_3^f)^2 + (I_3^f - 2\sin^2\theta_W Q_f - 2Q_f \cdot \Delta)^2} \tag{49}$$

where the correction

$$\Delta = \delta \sin^2\theta_W + \frac{\cos\theta_W}{\sin\theta_W} \cdot \frac{Re\hat{\Sigma}^{\gamma Z}(M_Z^2)}{M_Z^2} \tag{50}$$

[4]this may also be the as yet unknown "old" physics objects M_H, m_t

has two sources: first the shift in $\sin^2\theta_W$ due to virtual particles contributing to Δr; second the modification of the vector coupling by the $\gamma - Z$ mixing propagator:



Both contributions to Δ are universal for all types of external fermions. From (49) it follows that new physics in SU(2)×U(1) will contribute to A_f in terms of a shift

$$\delta A_f = \frac{4\,Q_f\,I_3^f\left[(I_3^f - 2\sin^2\theta_W Q_f)^2 - (I_3^f)^2\right]}{\left[(I_3^f - 2\sin^2\theta_W Q_f)^2 + (I_3^f)^2\right]^2}\cdot\Delta(new)$$

$$\equiv B(I_3^f, Q_f)\cdot\Delta(new) \qquad (51)$$

with a factor B independent of new physics and a term $\Delta(new)$ independent of the fermion species. (47) and (51) imply that the combination

$$\Delta_f = \delta\left[A_{LR} - \frac{4}{3}\cdot\frac{B(e)}{B(f)}\cdot A_{FB}^{pol}(e^+e^- \to f\bar{f})\right] \qquad (52)$$

vanishes in SU(2)×U(1) if others than propagator corrections are absent.

The presence of a new Z boson in the neutral current sector of (44) gives rise to deviations in A_{LR} and A_{FB}^{pol} which do not cancel in the difference (52) yielding $\Delta_f \neq 0$. Quantitatively the magnitude of Δ_f depends on the details of the model under consideration, in particular on the quantum numbers of the involved fermions with respect to the extra $U(1)_{Y'}$. Large deviations from $\Delta_f = 0$ hence can be interpreted as signals for a modication of the structure of the weak neutral current.

An example for $b\bar{b}$ and $c\bar{c}$ final states is shown in Figure 19. The fermion representations are those of the [27] of the E_6 unifying group; furthermore it is assumed that $g_1 = g_1'$. If we assume that A_{FB}^{pol} for b quark final states can be measured with an uncertainty of 0.02 Figure 19 indicates a sensitivity of Δ_f to a Z' with mass up to $M_{Z'} \approx 10 M_Z$.

The quantity Δ_μ for $e^+e^- \to \mu^+\mu^-$ is also a probe for the $e - \mu$ universality of the new Z': in case of universal couplings Δ_μ would be zero. We have already encountered another possible source for a non-universal coupling structure: non-standard Higgs bosons with enhanced Yukawa couplings (section 4.4). However, in view of the small μ mass which would require extremely different vacuum expectation values , it seems very unlikely that a violation of lepton universality is due to Higgs bosons.

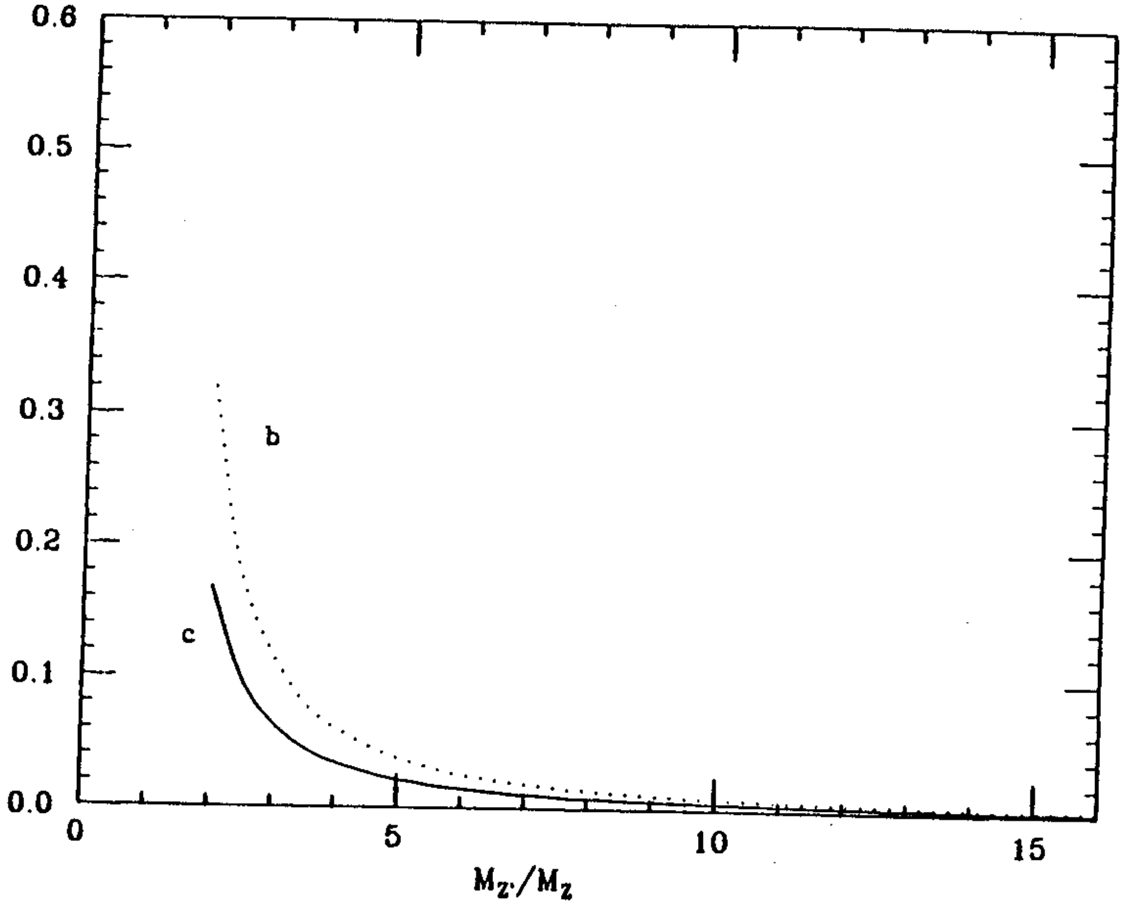


Figure 19: Δ_f, eq. (52), for $f = b$, c final states (from [20])

7 Summary

Forthcoming experiments at e^+e^- colliders will determine the vector boson masses and the on-resonance asymmetries A_{LR} and A_{FB}, A_{FB}^{pol} with high accuracy. In connection with the precisely measured μ decay constant G_μ these experiments provide precision tests of the standard model after the radiative corrections have been taken into account carefully. The largest part of the radiative corrections are in general the QED corrections, which demand a treatment beyond the one-loop level. For the determination of the Z resonance shape (measurements of mass and width) and for the left-right asymmetry these corrections are under control.

The high precision experiments will also become sensitive to possible new physics beyond the minimal model. This is of particular importance for new particles which are too heavy to be produced directly. As a powerful instrument longitudinal beam polarization is one of the keys to the presumably small new physics effects which are expected to be of the order of the conventional radiative corrections. As a measurable quantity with small systematic errors the left-right asymmetry is a sensitive probe for all additional particles within SU(2)×U(1) multiplets with isospin mass splittings as well as for additional vector bosons not within the minimal gauge group. It allows furthermore to determine small deviations from the standard ρ parameter value $\rho = 1$ which would be a signal for a Higgs sector with others than doublets and singlets.

Final state polarization or forward-backward asymmetries with polarized beams probe the universality of the weak fermion couplings, the existence of non-standard Higgs bosons with enhanced Yukawa couplings, and allow in combination with A_{LR} to isolate effects from extra neutral Z' bosons in extended electroweak symmetry groups. Presently the unknown top quark mass is the major source of uncertainty in revealing new physics effects. Theoretically the computation of complete radiative corrections for extended models would be desirable since the presently neglected terms might reach the level of the experimental accuracy.

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References:

1. Proceedings of the International Workshop on Radiative Corrections in SU(2)×U(1), Trieste 1983, ed. B.W. Lynn and J.F. Wheater, Singapore 1984

2. G Altarelli et al., in: Physics with LEP, eds. J Ellis and R Peccei, CERN 86-02

3. A. Barroso et al., in: Proceedings of the ECFA Workshop on Physics with LEP 200, Aachen 1986, ed. A. Böhm and W. Hoogland, CERN 87-08, ECFA 87-108

4. F.A. Berends, G.J.H. Burgers, W.L. van Neerven, Phys. Lett. 185 B (1987) 395; CERN-TH 4772/87

5. F.A. Berends, G.J.H. Burgers, W. Hollik, W.L. van Neerven, Phys. Lett. 203 B (1988) 177;

D. Yu. Bardin, A. Leike, T. Riemann, M. Sachwitz, Zeuthen preprint PHE 88-03 (1988);
V.S. Fadin, V.A. Khoze, Novosibirsk preprint 87-157 (1987)

6. B.W. Lynn, M. Peskin, R.G. Stuart, in: Physics with LEP, CERN 86-02

7. A. Sirlin, Phys. Rev. D 22 (1980) 971

8. M. Böhm, W. Hollik, H. Spiesberger, Fortschr. Phys. 34 (1987) 687

9. W.J. Marciano, Phys. Rev. D 20 (1979) 274;
W.J. Marciano, A. Sirlin, Phys. Rev. D 22 (1980) 2695

10. F. Jegerlehner, Z. Phys. C 32 (1986) 195

11. F. Jegerlehner, Z. Phys. C 32 (1986) 425

12. P. Roudeau et al., in: Proceedings of the ECFA Workshop on Physics with LEP 200, Aachen 1986, ed. A. Böhm and W. Hoogland, CERN 87-08, ECFA 87-108

13. M. Böhm, W. Hollik, Nucl. Phys. B 204 (1984) 45;
S. Jadach, J.H. Kühn, R.G. Stuart, preprint MPI-PAE/PTh 71/87;
J.H. Kühn, these proceedings

14. J. Jersak, E. Laerman, P.M. Zerwas, Phys. Rev. D 25 (1980) 1218;
R. Kleiss, F.M. Renard, C. Verzegnassi, Nucl. Phys. B 286 (1987) 669;
A. Djouadi, Montpellier preprint, PM/87-52 (1957)

15. A. Blondel, B.W. Lynn, F.M. Renard, C. Verzegnassi, Montpellier preprint PM/87-14 (1987)

16. Z. Wąs, these proceedings

17. J.C. Pati, A. Salam, Phys. Rev. D 10 (1974) 275;
R.N. Mohapatra, J.C. Pati, Phys. Rev. D 11 (1975) 566, 2554

18. V.G. Angelopoulos, J. Ellis, D.V Nanopoulos, N.D. Tracas, Phys. Lett. 176 B (1986) 203;
J. Ellis, K. Enqvist, D.V. Nanopoulos, F. Zwirner, Mod. Phys. Lett. A 1(1986) 57; Nucl. Phys. B 276 (1986) 14

19. P. Franzini, F. Gilman, Phys. Rev. D 32 (1985) 237; SLAC-PUB 3932 (1986);
G. Bélanger, S. Godfrey, Phys. Rev. D 34 (1986) 1309;
I. Bigi, M. Cvetič, Phys. Rev. D 34 (1986) 1651

20. M. Cvetič, B.W. Lynn, Phys. Rev. D 35 (1987) 42

21. M. Veltman, Acta Phys. Pol. B 8 (1977) 475

22. H.E. Haber, G.L. Kane, Phys. Rep. 117 (1985) 75

23. J.A. Grifols, J. Solà, Nucl. Phys. B 253 (1985) 47

24. CHARM II Proposal, C. Busi et al., CERN/SPSC/83-24 (1983)

25. H.P. Nilles, Phys. Rep. 110 (1984) 1;
J. Ellis, in: "Superstrings and Supergravity", eds. A.T. Davis, and D.G. Gutherland, SUSSP publications, Edinburgh 1986

26. J. Ellis, S. Rudaz, N.D. Tracas, CERN-TH.4657/87 (1987)

27. A.J. Buras, Proc. of the EPS Conference, Bari 1985, and references therein

28. R.D. Peccei, H.R. Quinn, Phys. Rev. Lett. 38 (1977) 1440; Phys. Rev. D 16 (1977) 1719

29. A. Axelrod, Nucl. Phys. B 209 (1982) 349;
M. Clements, C. Footman, A. Kronfeld, S. Navasimhan, D. Photiadis, Phys. Rev. D 27 (1983) 570;
V. Ganapathi, T. Weiler, E. Laerman, I. Schmitt, P.M. Zerwas, Phys. Rev. D 27 (1983) 579

30. S. Bertolini, Nucl. Phys. B 272 (1986) 77

31. W. Hollik, Z. Phys. C 32 (1986) 291; Z. Phys. C 37 (1988) 569

32. J.F. Gunion, H.E. Haber, Nucl. Phys. B 272 (1986) 1

33. CELLO Collaboration, DESY 87-030 (1987)

34. H.E. Haber, G.L. Kane, T. Sterling, Nucl. Phys. B 161 (1979) 493

35. A.A. Akhundov, D. Yu. Bardin, T. Riemann, Nucl. Phys. B 276 (1986) 1;
W. Beenakker, W. Hollik, DESY 88-007 (1988), Z. Phys. C (to appear)

36. J.L. Hewett, S. Nandi, T.G. Rizzo, Ames Laboratory preprint OSU 201, IS-J-2983 (1988)

37. U. Amaldi et al., Phys. Rev. D 36 (1987) 1385;
G. Costa et al., CERN-TH.4675/87 (1987)

38. B.W. Lee, C. Quigg, H.B. Thacker, Phys. Rev. D 16 (1977) 1519;
M. Chanowitz, M. Golden, H. Georgi, Phys. Rev. Lett. 57 (1986) 2344;
Phys. Rev. D 36 (1987) 1490

39. U. Baur, M. Lindner, K.H. Schwarzer, preprint MPI-PAE/PTh 74/86, FERMILAB-PUB-86/136-T (1986)

40. R. Casalbuoni, P. Chiappetta, D. Dominici, F. Feruglio, R. Gatto, CERN-TH.4876/87 (1987)

41. W. Hollik, Z. Phys. C 8 (1981) 149;
W. Hollik, A. Zepeda, Z. Phys. C 12 (1982) 67

42. B.W. Lynn, F.M. Renard, C. Verzegnassi, Montpellier preprint PM-87/45 (1987); F. Boudjema, F.M. Renard, C. Verzegnassi, preprint "New Physics Sum Rules from Polarized Electron Positron Asymmetries on Z Resonance"; C. Verzegnassi, these proceedings

43. J. van der Bij, M. Veltman, Nucl. Phys. B 231 (1984) 205