SEARCH FOR THE HIGGS AT FUTURE ACCELERATORS

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1 Introduction

The experimental verification of the Standard Model, which is made up of QCD [1] and the electroweak theory [2], is still to be completed in that the top quark and the Higgs boson [3] have not yet been found. Both are important. The large mass of the t quark as compared with the other known quarks might indicate that its couplings are perhaps different from the Standard Model prediction (e.g. by mixing with some heavier exotic state [4]). Also the knowledge of the t mass is essential to sharpen the Standard Model predictions, thus allowing more stringent precision tests of the theory at LEP and elsewhere. However, there is no doubt that the most essential problem facing experimental particle physics in the next decade is the question of the physical origin of the electroweak symmetry breaking.

If the Standard Model is a reliable guidance, the top quark should be found in the next few years at the Tevatron of Fermilab. In fact, assuming the Standard Model, from precision tests of the electroweak theory the limit [5] $m_t \leq 200$ GeV is derived and, actually, values around $m_t \simeq 130$–140 GeV are favoured. The search for the Higgs is being pursued at LEP 1 and will continue at LEP 200. Indeed all previous limits on the Higgs mass m_{H} have been dwarfed by only a few months of LEP operation. For the standard Higgs we have at present the following results [6]:

$$\begin{aligned} \text{ALEPH} &: \; m_{\mathrm{H}} > 48 \text{ GeV} \\ \text{OPAL} &: \; m_{\mathrm{H}} > 44 \text{ GeV} \\ \text{L3} &: \; m_{\mathrm{H}} > 41.8 \text{ GeV} \\ \text{DELPHI} &: \; m_{\mathrm{H}} > 41 \text{ GeV} \end{aligned} \tag{1.1}$$

These limits are obtained from negative searches of the process $e^+e^- \to HZ^* \; (\to f\bar{f})$.

As is well known, the value of the Higgs mass is not predictable even in the minimal Standard Model with a single Higgs doublet. What is certainly true is that the Higgs boson cannot be too heavy or the perturbative theory becomes sick and breaks down [7]. If $m_{\mathrm{H}} \geq O(1 \text{ TeV})$ the perturbative rates for VV $\to$ VV scattering (V $\equiv$ W, Z) violate the unitarity limit [8] for $\sqrt{s} \gg m_{\mathrm{W}}$. More important than this, in non-asymptotically free gauge theories there are Landau poles where the coupling constant blows up according to the renormalization group improved perturbation theory (unless the renormalized coupling is not vanishing so that the theory is a free theory, i.e. trivial). This phenomenon is also present in QED but it would only occur beyond the Planck scale of mass, so that the problem can be solved at such large energies by embedding the theory in a larger context (e.g. grand unification). The coupling of the quartic term $\lambda(\phi^+\phi)^2$ in the Higgs potential increases with m_{H}^2 ($m_{\mathrm{H}}^2 \sim \lambda/G_{\mathrm{F}}$). In addition, for a given m_{H}, λ increases logarithmically with energy because the theory is not asymptotically free in the Higgs sector. Thus the position of the Landau pole depends on m_{H}. Imposing that the Landau pole is far enough for the theory to make sense up to a scale Λ, gives a bound [9] on the standard Higgs mass which is plotted in Fig. 1, taken from Ref. [10]. We see that for a

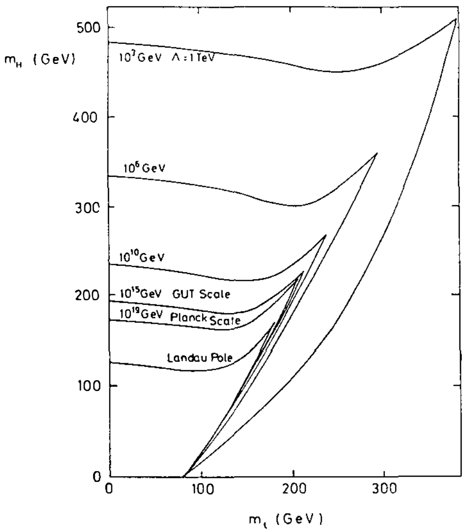

Fig. 1 Combined limits (from Ref. [10]) on m_H and m_t from vacuum stability and avoiding the Landau pole up to a scale Λ.

light Higgs, i.e. $m_H \leq 180$–200 GeV, the perturbative regime is valid up to M_{GUT} or M_{Pl}. For a heavier Higgs the value of Λ decreases until eventually $\Lambda \sim m_H$. For $m_H \sim 1$ TeV, the theory is valid up to $\Lambda \sim 1$ TeV.

We can understand these results by the following crude simplification [7]. The renormalization group equation for the quartic coupling λ, in the limit of neglecting gauge and Yukawa couplings, becomes:

$$\frac{d\lambda(t)}{dt} = \frac{3}{4\pi^2} \lambda^2(t) , \tag{1.2}$$

with $t = \ln \Lambda/v$, where v is the Higgs vacuum expectation value and Λ is the scale where λ is evaluated. The coefficient $\beta_0 = 3/4\pi^2$ is obtained from one-loop corrections to the quartic coupling in the $\lambda(\phi^+\phi)^2$ theory. The normalization of v and λ, in physical terms, is here chosen such that

$$\lambda \equiv \lambda(v) = \sqrt{2}\, G_F\, m_H^2 \tag{1.3}$$

$$v = (2\sqrt{2}\, G_F)^{-1/2} \simeq 174 \text{ GeV} . \tag{1.4}$$

When $\lambda(t)$ is large the gauge and Yukawa couplings can be neglected with the exception of the top Yukawa coupling, which can become large if $m_t \geq v : g_{top} = m_t/v$. By solving (1.2) one obtains

$$\lambda(t) = \frac{\lambda}{1 - 3/4\pi^2 \lambda t} . \tag{1.5}$$

The minus sign in the denominator, typical of non-asymptotically free theories, implies the increase of $\lambda(t)$ with the scale Λ up to an infinite value which is obtained for $3/4\pi^2\ \lambda t = 1$. In order to avoid the Landau pole the condition

$$\frac{3}{4\pi^2}\ \lambda t \ = \ \frac{3}{4\pi^2}\sqrt{2}\, G_\mathrm{F}\, m_\mathrm{H}^2\ \ln\ \Lambda/v < 1 \tag{1.6}$$

must be imposed. This condition is equivalent to:

$$\frac{m_\mathrm{H} \leq 893 \text{ GeV}}{\sqrt{\ln\ \Lambda/174 \text{ GeV}}} \tag{1.7}$$

or $m_\mathrm{H} < 144,\ 165,\ 675$ GeV for $\Lambda = 10^{19}, 10^{15}, 10^{3}$ GeV respectively.

We see that this simple model reproduces the quantitative features of the bounds on m_H in Fig. 1 fairly well. The curves in Fig. 1 are obtained by a more refined renormalization group treatment of the problem, with inclusion of gauge and top effects. The obvious criticism to the above approach is that a perturbative evaluation of the β function is not justified in the vicinity of the Landau pole. Thus it is very interesting that the validity of the bound has been confirmed by recent computer simulations of the electroweak theory on the lattice [11]. The precise value of the upper limit on m_H depends on the exact definition of Λ and on where one fixes the line between acceptable and not acceptable. In fact the lattice results nicely extrapolate the perturbative evaluation (Fig. 2) and find limits on m_H such that:

$$m_\mathrm{H} \leq (8\text{–}10) m_\mathrm{W} \simeq 0.6\text{–}0.8 \text{ TeV}\ . \tag{1.8}$$

It is thus fair to conclude that the internal consistency of the Standard Model demands that the Higgs mass is below 1 TeV.

In Fig. 1 there is also a forbidden region at large m_t and small m_H. This boundary is determined by the requirement of vacuum stability [7, 12].

At tree level the scalar potential is given by

$$V(\varphi) = -\mu^2|\varphi|^2 + \frac{\mu^2}{2v^2}|\varphi|^4\ . \tag{1.9}$$

The quantum corrections can be computed by expanding in the number of loops. At one loop one obtains:

$$V(\varphi) = -\mu^2|\varphi|^2 + \frac{\mu^2}{2v^2}|\varphi|^4 + \gamma|\varphi|^4\left(\ln\frac{|\varphi|^2}{v^2} - \frac{1}{2}\right) \tag{1.10}$$

with

$$\gamma = \frac{3\sum\limits_{\text{vectors}} m_\mathrm{v}^4 + \sum\limits_{\text{scalars}} m_\mathrm{s}^4 - 4\sum\limits_{\text{fermions}} m_\mathrm{f}^4}{64\pi^2 v^4} \tag{1.11}$$

It is simple to check that also in the corrected form v is an extremum of $V(\phi)$. In the minimal Standard Model with one Higgs doublet and three fermion families one obtains

$$\gamma = \frac{3m_Z^4 + 6m_W^4 + m_H^4 - 12m_t^4}{64\pi^2 v^4} . \tag{1.12}$$

The extra factor of three in front of m_t^4 is of course due to colour.

For the realization of spontaneous symmetry breaking and stability of the theory one requires that i) the extremum at $\varphi = v$ is a minimum, i.e. $V(v) < 0$ and ii) $V(\varphi) \to +\infty$ for $|\varphi| \to \infty$, so that the Hamiltonian is bound from below.

At small m_t, $m_t < 80$ GeV, the first requirement leads to the Linde–Weinberg limit $m_H^2 > 2\gamma v^2$, or $m_H \geq 7$ GeV. This limit is by now void, because of the experimental lower bounds on both m_t and m_H. For the second requirement to be fulfilled, m_H must increase with m_t in order to prevent γ from becoming too negative [12]. At large $|\varphi|$ the one-loop evaluation of $V(\varphi)$ is not sufficient, and one needs a resummation of the large logarithms $\log \varphi^2/v^2$. The results are shown in Fig. 1 [10]. The above limits are only valid in the minimal Standard Model with one Higgs doublet. Note that in case that there are two or more Higgs doublets the limits refer to some average mass. Thus for the lightest Higgs the lower limit can be easily evaded but the upper limit is *a fortiori* valid. In conclusion either the Higgs is found below $\simeq 1$ TeV or new physics beyond the Standard Model should appear. At least one should see the onset of a new non-perturbative regime where the weak interactions become strong.

There is a widespread opinion among theorists that there must be some new physics beyond the Standard Model at a scale of energy of O(1 TeV). It is considered implausible that the origin of the electroweak symmetry breaking can be explained by the standard Higgs mechanism without accompanying new physics. The argument is one of naturalness and runs along the following lines. In the $SU(2) \otimes U(1)$ symmetric limit there are no masses. Both the gauge bosons and the fermions are massless. After symmetry breaking, all masses are proportional to the Higgs vacuum expectation value v or equivalently to $G_F^{-1/2} \simeq 293$ GeV ($v = 2^{-3/4} G_F^{-1/2}$) which is called the weak (or Fermi) scale. This is the characteristic scale of the electroweak theory. While the smallness of the Yukawa couplings that determine the light fermion masses and their ratios is not understood (but this problem can perhaps be solved at the level of the theory at very large energy scales), it remains true that $G_F^{-1/2}$ is the scale of mass of the theory. As is well known, a direct extrapolation of the Standard Model leads to grand unified theories [13] at a scale $M_{GUT} \simeq 10^{14}$–10^{16} GeV, close to the scale of quantum gravity $M_{Pl} \simeq 10^{19}$ GeV. One is perhaps led to imagine a unified theory of all interactions, including gravity (at present the best attempt at such a theory is provided by superstrings [14]). But certainly particle physics can no longer ignore such large scales of mass as M_{GUT} and M_{Pl}. Indeed, going from $G_F^{-1/2}$ up to M_{Pl} is an enormous gap of about 17 orders of magnitude. The obvious question is whether the Standard Model can extend its validity up to M_{Pl}. The answer is that this appears unlikely (the hierarchy problem). A natural explanation of $M_{Pl}/G_F^{-1/2} \simeq 10^{17}$ demands the presence of new physics near $G_F^{-1/2}$. The reason is that, if the Standard Model is valid up to a large scale Λ, even if one sets a small value for m_H at the tree (classical) level, $m_H \ll \Lambda$, the loop (quantum) corrections would make m_H increase up to the order of Λ.

The problem is especially acute for scalar fields because the corresponding mass divergences are quadratic, while they are only logarithmic for spin 1/2 fermions. Note that here the discussion is on the relation between bare and renormalized masses, where the cut-off dependence is hidden. In the renormalization procedure a physical value is simply assigned to m_{H} and it is left to the bare mass and the cut-off to adjust to each other. The naturalness problem arises if the divergences are seen as a low-energy effect, to be eventually removed by some new physics at the scale Λ (e.g. by gravity at M_{Pl}). Then the large momentum cut-off and the scale of new physics can be physically identified. The quadratic divergences associated with scalars are unacceptable in a 'natural' theory, while the logarithmic singularities of fermion masses can be tolerated. The fermion masses are also protected by chiral symmetry, which demands mass corrections to vanish in the massless limit, i.e. $\delta m \sim m \ln \Lambda/m$.

One possible solution is that the Higgs doublet really consists of fundamental scalar fields but naturality is restored by broken supersymmetry [15]. In the supersymmetric limit there is complete boson–fermion symmetry. The quadratic mass divergences associated with scalars cancel away so that only logarithmic singularities for both scalars and fermions are present. When supersymmetry breaking is switched on, the scale for δm_{H}^2 is naturally set by the splitting between partners in supersymmetric multiplets. The Fermi scale is natural if the masses of sparticles are around the Fermi scale. In the limit of exact supersymmetry and exact gauge $\mathrm{SU}(2)\otimes\mathrm{U}(1)$, all particles are massless. When supersymmetry is broken while $\mathrm{SU}(2)\otimes\mathrm{U}(1)$ is still preserved, ordinary particles remain massless while sparticles become massive. It is important to note that observed particles are precisely those whose mass terms are forbidden in the $\mathrm{SU}(2)\otimes\mathrm{U}(1)$ limit, while sparticle masses are allowed. For example, quark and lepton masses are forbidden while squark or slepton masses are allowed, the gauge boson masses are forbidden but the gaugino masses are allowed. Thus the fact that all ordinary particles were observed but no sparticles is not unnatural. When finally the $\mathrm{SU}(2)\otimes\mathrm{U}(1)$ symmetry breaking is switched on, the scalar mass naturally takes a value of the order of the scale of sparticle masses and all ordinary particles acquire a mass.

Many theorists working on quantum gravity and superstrings tend to consider SUSY as 'established' at M_{Pl} and beyond. For economy one is then naturally led to try to use SUSY at low energy, in order to solve some of the problems of the Standard Model. It is thus very important that it was indeed shown [16] that models where SUSY is softly broken by gravity do offer a viable alternative. The minimal supersymmetric Standard Model [17] (MSSM), which will be often mentioned in the following, is a well-specified theory, completely consistent and, in some respects, better than the Standard Model, as we have seen. The supersymmetric option is very appealing to theorists. It would represent the ultimate step of a continuous line of progress obtained by constructing field theories with an increasing degree of exact and/or broken symmetry and applying them to fundamental interactions. The value of the ratio of knowledge versus ignorance would be remarkably large in the case of SUSY: the correct degrees of freedom for a description of physics up to gravity would have been identified, the Hamiltonian would be known—apart from the values of a number of parameters— and the theory would be, to a large extent, computable up to the Planck scale.

The alternative main avenue to solve the hierarchy problem is to avoid fundamental scalar fields at all. This necessarily implies the existence of new strong forces. For example, the electroweak symmetry could be broken by condensates of new fermions attracted by a new force with $\Lambda_{new} \simeq G_F^{-1/2}$, Λ_{new} being the analogue of Λ_{QCD}, as in technicolour theories [18]. The mechanism that gives mass to $W^{\pm}$ and Z would be the analogue for a gauge theory of the breaking of chiral symmetry (a global symmetry) in QCD. A new anomaly-free multiplet of heavy technifermions, bound by a very strong gauge force called technicolour, must be introduced. The longitudinal modes of $W^{\pm}$, Z would be analogous to the pions in QCD. This approach faces problems [19] related to the existence of additional light pseudo-Goldstone bosons that should have been detected. In addition the fermion masses remain an unsolved question (the so-called extended technicolour introduced [20] to solve this problem leads more to new difficulties than to advantages).

Recently it has been proposed [21] that a very heavy top mass ($m_t \geq 230$ GeV) could induce the electroweak symmetry breaking. The Higgs would be a sort of $t\bar{t}$ bound state with mass $m_H \simeq 1.1 m_t$. This model (of the Nambu–Jona Lasinio type [22]) is non-renormalizable and involves many *ad hoc* four-fermion interactions to fix the fermion masses. More generally, the Higgs could be a composite of new fermions bound by a new force [23]. Or the $SU(2) \otimes U(1)$ symmetry could be a low-energy fake [24]. At large energies, $E \gg G_F^{-1/2}$, the $W^{\pm}$ and Z would be resolved into their constituents.

However it is fair to say that the above ideas become increasingly generic (going down the list). No sound theoretical framework has been developed out of them. The compositeness alternative, in all its different forms, is not at all so neatly formulated as the supersymmetric option. On the contrary, in many respects the compositeness way is not well defined at all and leads to many unsolved problems. But, of course, this state of affairs could only be due to a lack of ingenuity on the part of theorists.

In conclusion there are solid arguments for new physics near the Fermi scale of mass $G_F^{-1/2}$. Either a fundamental scalar Higgs exists and naturalness is restored by supersymmetry, or new strong forces will manifest themselves, drastically changing the framework of the Standard Model beyond O(1 TeV). A new non-perturbative regime will set up, with new resonances, and the physics will become less predictable above that energy. An important point is that all conceivable possibilities are very complex. Each of them implies a rich new spectrum of states and phenomena: the whole spectrum of superpartners in SUSY; new hadrons, excited vector bosons, etc., in the composite alternative. The new physics is in all cases distributed over a large interval of energies. The low-lying fringes of the new spectroscopy, or at least their virtual effects, should already be accessible to LEP 1 and LEP 200. A lot of discoveries are expected at the LHC, to be followed by more at the SSC.

2 Search for the Standard Higgs

It is clear that no other accelerator is better than LEP for finding a Standard Model Higgs with mass $m_H \leq m_Z$. We have already mentioned the present lower limits on m_H obtained at

LEP 1 [6]. In the next few years of continued LEP 1 operation one can presumably improve the limits up to $m_H \leq 60$ GeV. Beyond that the increase of energy is absolutely necessary. At LEP 200, which will be operational at the end of 1993, the range m_H = 50–90 GeV can be explored. The LEP 200 process for observing the Higgs is $e^+ + e^- \rightarrow ZH$ (with a real Z in the final state and not a virtual one as in the analogous process at LEP 1). The observation of the Standard Model Higgs with $m_H \leq 80$ GeV is considered an easy problem at LEP 200 with project energy and luminosity [25]. For $m_H \sim$ 80–95 GeV the problem is considerably more difficult because of the small cross-section and of the H/Z confusion due to the overlapping of masses. This case was studied recently by Kunszt and Stirling [26]. The total cross-sections for the signal ($e^+e^- \rightarrow HZ$) and the main background ($e^+e^- \rightarrow ZZ$) are shown in Fig. 2 for $\sqrt{s}$ = 200 GeV. The signal cross-section is small ($\sim$ 0.5 pb without branching ratios).

In the channel $\ell^+\ell^-$ + jets, ℓ = e, μ, with $\int L\, dt$ = 500 pb^{-1}, the ratio of the numbers of events for signal and background [26] is 20.4/2.0, 17.3/28, 15.1/5.1 for m_H = 85, 91.1, 95 GeV respectively, with 'm_Z' = $m_H \pm 1$ GeV, where 'm_Z' is the reconstructed mass of $Z \rightarrow$ jets. A moderate help can be obtained from cuts in $\cos\theta$ (θ = $\mu^{\pm}$–beam angle) as the angular distributions are different. The conclusion of Ref. [26] is that, close to m_Z, high energy ($\sqrt{s} \simeq 200$ GeV) and large luminosities ($\int L\, dt \geq 500$ pb^{-1}) are needed.

Beyond LEP 200 the future of e^+e^- accelerators is probably in linear colliders. The search for the Higgs at linear e^+e^- colliders with $\sqrt{s}$ = 1–2 TeV was discussed at the La Thuile Workshop [27] and elsewhere. At present an international Workshop is being organized by ICFA to study the physics potential of an e^+e^- collider with $\sqrt{s}$ = 0.5 TeV and $L = 10^{33}$ cm^{-2} s^{-1}. The results will be presented in Finland next September. For the intermediate-mass Higgses with

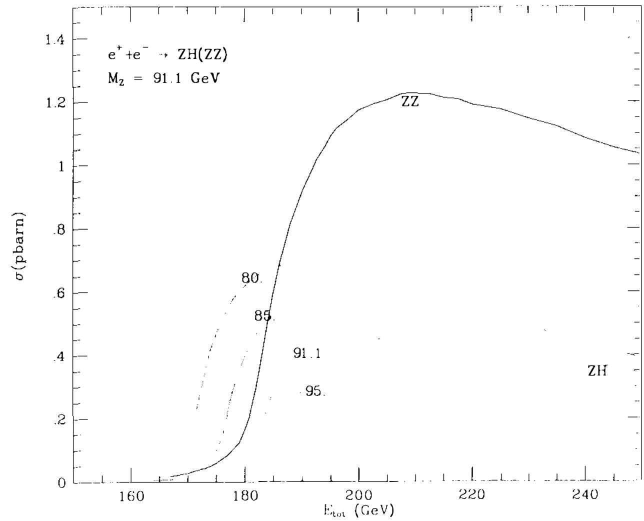

Fig. 2 Total cross-sections for $e^+e^- \rightarrow ZZ$ and $e^+e^- \rightarrow ZH$ for m_H = 80, 85, 90, and 95 GeV as functions of the e^+e^- c.m. energy (from Ref. [26]).

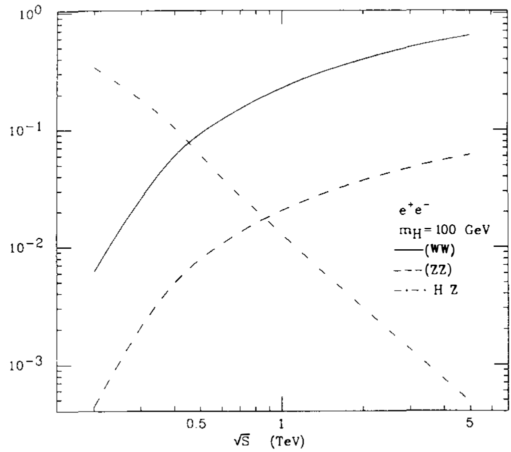

Fig. 3 Cross-section for $e^+e^- \rightarrow H + ...$ as a function of energy for different mechanisms (Ref. [27]); WW and ZZ indicate the corresponding fusion channels (e.g. WW stands for $e^+e^- \rightarrow H\nu\bar{\nu}$), while ZH means $e^+e^- \rightarrow$ ZH (via Z exchange).

$m_Z \leq m_H \leq 2m_Z$ (the region which is difficult at the LHC and SSC) e^+e^- linear colliders are good. It turns out (Fig. 3) that $\sqrt{s} = 0.5$ TeV is just the energy where the cross-sections of $e^+e^- \rightarrow$ HZ and $e^+e^- \rightarrow \nu\bar{\nu}$H (via W–W fusion) are equal for $m_H \simeq 100$ GeV [27]. For m_H in this region, $\geq 10^3$ events/year can be expected at $\sqrt{s} \geq 0.5$ TeV with $L = 10^{33}$ cm^{-2} s^{-1}. The γ–γ background can be controlled [28].

The search for the minimal Standard Model Higgs [29] at the LHC and the SSC has been discussed in great detail at the Aachen Workshop [30] as well as at previous ones on LHC [27, 31] and SSC [32] physics. This is a good reference problem, but not necessarily the central issue of physics at the LHC. After all the Higgs might be found at LEP. Such a discovery there would not at all mean that the LHC is no longer necessary. In fact, we have seen that one expects some new physics at the weak scale to accompany the Higgs. The minimal Standard Model might well be wrong for the Higgs sector. For example, the Higgs sector of supersymmetric models involves at least two Higgs doublets [17, 29]. The couplings of the lightest SUSY Higgs are not as in the minimal Standard Model. However, it would in many cases be impossible to prove at LEP that the Higgs candidate is the particle predicted by the minimal Standard Model. The Higgs search is a good reference problem in the sense that experiments must be good enough to see the standard Higgs in order to prove adequate for the solution of the electroweak symmetry-breaking question. The discovery of the Higgs is in fact a very difficult experimental problem, because the Higgs is heavy and, its couplings being proportional to masses, it is essentially not coupled to light particles (the most common ones). Heavy real or virtual states must be excited in order to produce the Higgs, so that the cross-sections are relatively small. In addition, below the WW or ZZ threshold, the dominant decay into the heaviest accessible pair of quarks is swamped by the QCD background. The case of the Standard Model Higgs was studied by a dedicated group at the Aachen Workshop, convened by Z. Kunszt and J. Stirling.

The problem was restarted from scratch. Calculations of the total width (Fig. 4) and of the branching ratios (Fig. 5) were updated by Z. Kunszt and J. Stirling. The inclusion of the effects from the running of the b-quark mass makes the $b\bar{b}$ partial width smaller, and the rare decay branching ratios below the $t\bar{t}$, WW, and ZZ thresholds larger. In particular the $H \to \gamma\gamma$ branching ratio was found to be larger by a welcome factor of 2 with respect to previous calculations. The production occurs mainly through gluon–gluon fusion ($gg \to H$) via a quark loop (dominated by virtual t exchange) or through WW fusion plus a small ZZ contribution ($q\bar{q} \to Hq\bar{q}$). For $m_t \geq 90$ GeV the gluon-fusion process is dominant up to very heavy Higgs masses: $m_H \geq$ 600 GeV for $m_t \simeq 90$ GeV, or $m_H \geq 1$ TeV for $m_t \simeq 180$ GeV (Fig. 6).

The intermediate-mass Higgs is the most difficult case. It is assumed that a light Higgs with mass $m_H \leq m_Z$ will be discovered at LEP 1 or LEP 200. The intermediate Higgs range is defined by $m_Z \leq m_H \leq 2m_Z$, i.e. below the threshold for $H \to ZZ$. This region would be hopeless if $H \to t\bar{t}$ were allowed. Now it is known from CDF results that indeed $m_t \geq m_Z$, so that the dominant decay of the intermediate Higgs is $H \to b\bar{b}$. This implies that the accessible decay modes $H \to ZZ^* \to 4\ell^\pm$ and $H \to \gamma\gamma$ have a much larger branching ratio. High luminosity, $L \simeq 10^{34}$ cm^{-2} s^{-1}, is absolutely necessary for detecting the intermediate Higgs at the LHC.

The first very important conclusion which was obtained is that with $\int L\, dt \simeq 10^5$ pb^{-1} and both e and μ detection, it is possible to observe the intermediate Higgs for $m_H \geq 130$ GeV through the chain $H \to ZZ^* \to 4\ell^\pm$ ($\ell = e, \mu$) [33]. The signal rate before cuts is 100–700 events

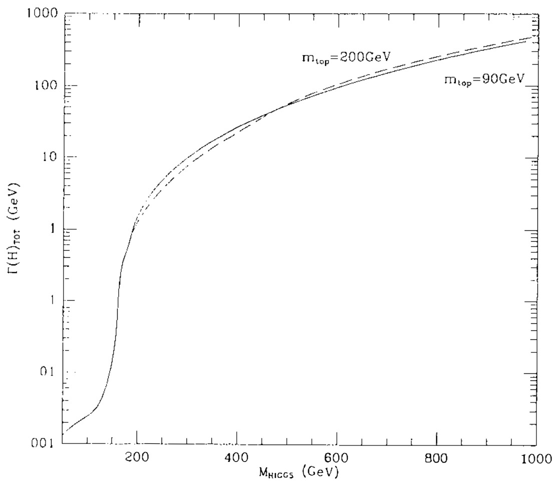

Fig. 4 The total width (from Ref. [30]) of the standard Higgs as a function of m_H (and m_t).

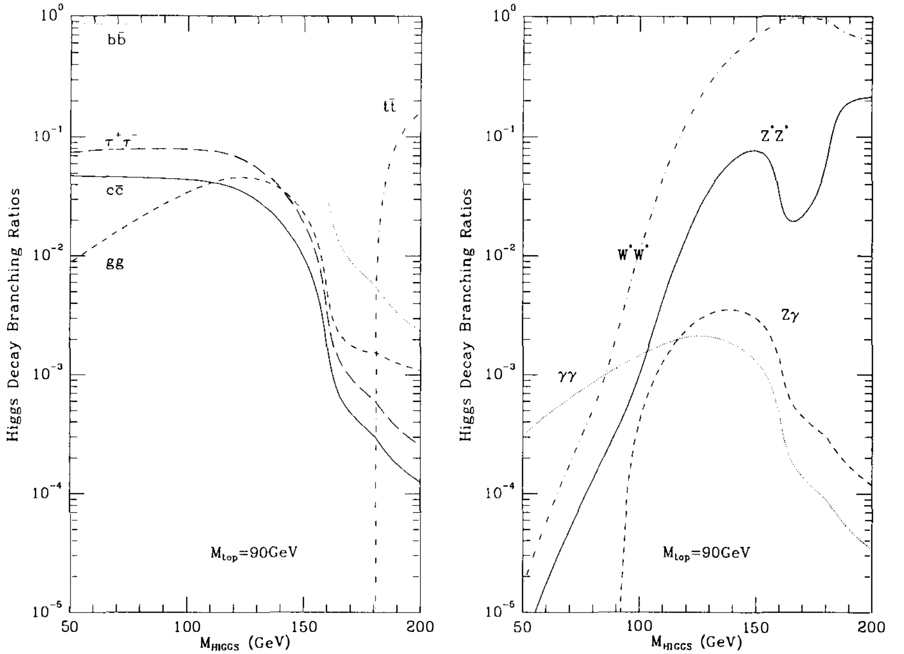

Fig. 5 The branching ratios (from Ref. [30]) of the standard Higgs. For $m_H > 200$ GeV, the WW and ZZ channels are dominant. There is little dependence on m_t.

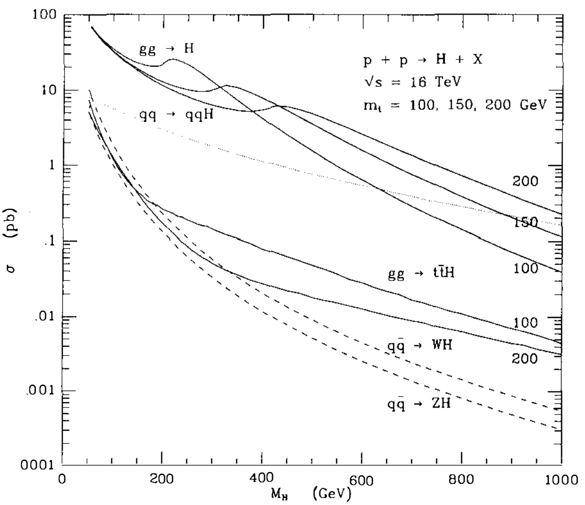

Fig. 6 Production cross-sections (from Ref. [30]) of the standard Higgs at the LHC.

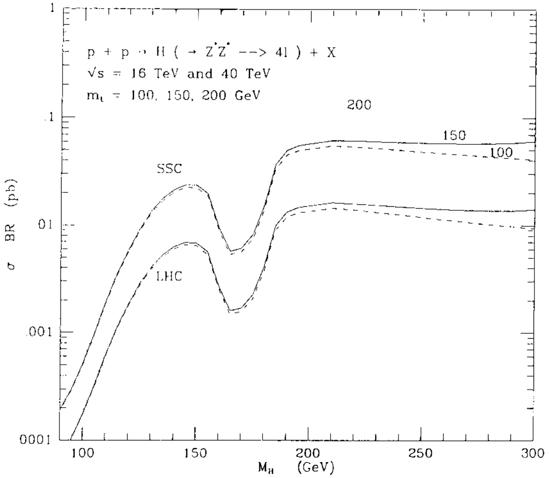

Fig. 7 The cross-section times branching ratio (from Ref. [30]) for pp $\rightarrow$ H($\rightarrow$ Z*Z* $\rightarrow 4\ell^{\pm}$)X at the LHC and SSC.

per year as seen from Fig. 7 (the dip at $m_{\rm H} \sim$ 160 GeV corresponds to the opening of the threshold for WW decay, which is not practicable because of the t$\bar{\rm t}$ $\rightarrow$ WWb$\bar{\rm b}$ background). A thorough study of backgrounds was done. Particular attention was devoted to the Zb$\bar{\rm b}$ channel (the leptons from b$\bar{\rm b}$ can be hard and isolated enough to mimic the Z*). The dominant process gg $\rightarrow$ Zb$\bar{\rm b}$ was studied by van Eijk and Kleiss in Ref. [30]. Detailed simulations of the t$\bar{\rm t}$, Zb$\bar{\rm b}$, Z*Z*, and Z$^*\gamma^*$ backgrounds were performed. The signal is already visible over the background without isolation cuts (Fig. 8a), but becomes much more prominent with isolation cuts (Fig. 8b).

Much work was devoted to the problem of closing the window $m_{\rm Z} \leq m_{\rm H} \leq 130$ GeV. This is a particularly hard task. The main line of attack is based on the process pp $\rightarrow$ H($\rightarrow \gamma\gamma$)X, first discussed in Ref. [33] and then widely studied [27, 32]. This process was further analysed at the present Workshop. I refer the reader to the article by C. Seez et al. for a detailed discussion [30]. The conclusion was that this channel is extremely difficul, but feasible with a very good detector. The signal rate is 0.5–1 $\times$ 10^3 events per $\int L\ {\rm d}t \simeq 10^5$ pb^{-1} (Fig. 9). The intrinsic background from q$\bar{\rm q}$ $\rightarrow \gamma\gamma$ and gg $\rightarrow \gamma\gamma$ already poses a formidable problem. A superb electromagnetic calorimeter is required, and vertex localization is very important for the $\gamma\gamma$ invariant-mass reconstruction. In Table 1 we show the comparison of signal versus intrinsic background for $m_{\rm H}$ = 80–150 GeV and $\int L\ {\rm d}t \simeq 10^5$ pb^{-1}.

The reducible background from jets misidentified as photons demands a large rejection factor $r_{2\rm j} = r_{1\rm j}^2 > 10^8$, where $r_{2\rm j}$ and $r_{1\rm j}$ correspond to double- and simple-jet misidentification, respectively. The possibility of a position detector, located some 2 m away, in order to see the separation between the two γ's from π^0 decays, was suggested as a main device for the discrimination of the jet background.

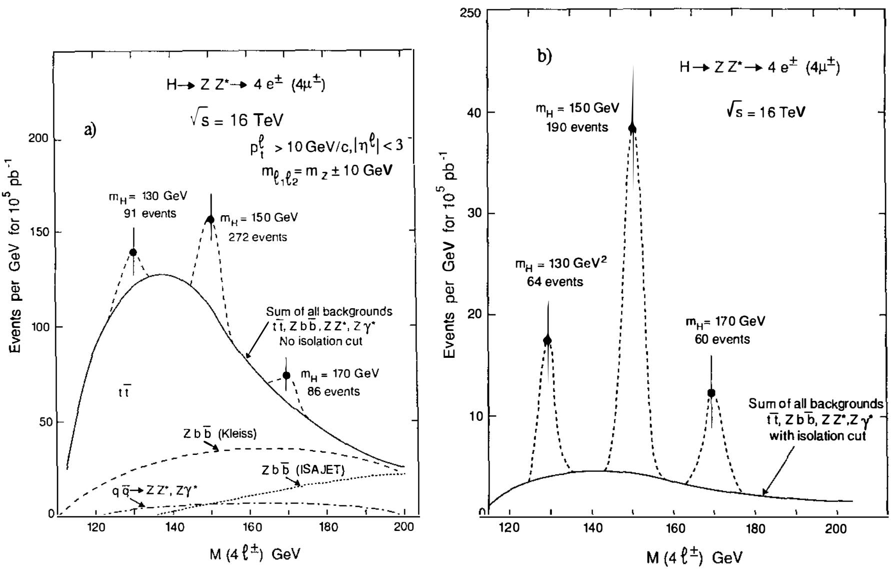

Fig. 8 The signal and background for the intermediate-mass Higgs ($H \to ZZ^* \to 4\ell^\pm$, $\ell = e, \mu$), a) without and b) with isolation cuts (from Ref. [30]).

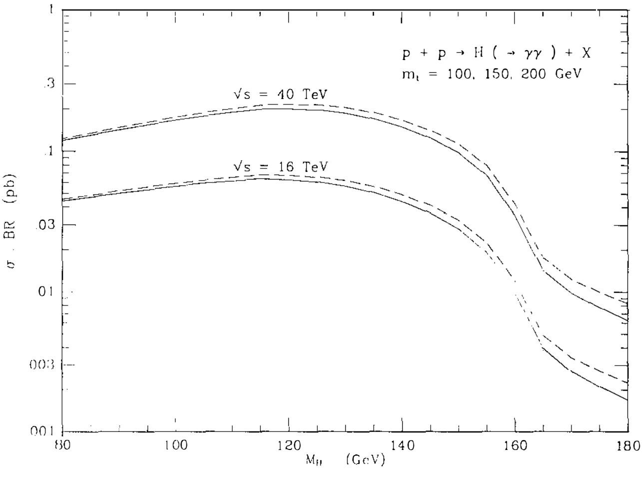

Fig. 9 The signal for $H \to \gamma\gamma$ at the LHC and SSC (from Ref. [30]).

Table 1 pp $\to$ H($\to \gamma\gamma$)X

$m_{\rm H}$ (GeV)	ΔM (GeV)	Signal	Background	S/$\sqrt{\rm B}$
80	1.0	570	11800	5.2
100	1.5	1180	13700	10.1
150	2.0	830	5600	11.1

An additional possibility, at small $m_{\rm H}$, is provided by the associated production of HW followed by H $\to \gamma\gamma$: pp $\to$ H($\to \gamma\gamma$)W($\ell\nu$)X. This process was studied at the Workshop by Kleiss, Kunszt, and Stirling [34], and, from the experimental point of view, by Di Lella et al. [30]. The good thing about this process is that the sum of the irreducible background from W$\gamma\gamma$ and of the reducible one from b$\bar{\rm b}$g, b$\bar{\rm b}\gamma$, b$\bar{\rm b}\gamma\gamma$, Wjj,..., with misidentifications, is very small in comparison with the signal. The bad thing is that the signal rate is also very small (Fig. 10) [34]. The resulting number of events for signal and background after cuts are collected in Table 2. It is concluded [30] that this channel is very difficult but could provide a useful way of confirming the signal from pp $\to$ H($\to \gamma\gamma$)X.

Table 2 pp$\to$ H($\to \gamma\gamma$)W($\to \ell\nu$)X

$m_{\rm H}$ (GeV)	Signal	Background		
		Irreducible	Reducible	Total
75	17	6	1	7
100	22	3	1	4
130	18	2	< 1	3

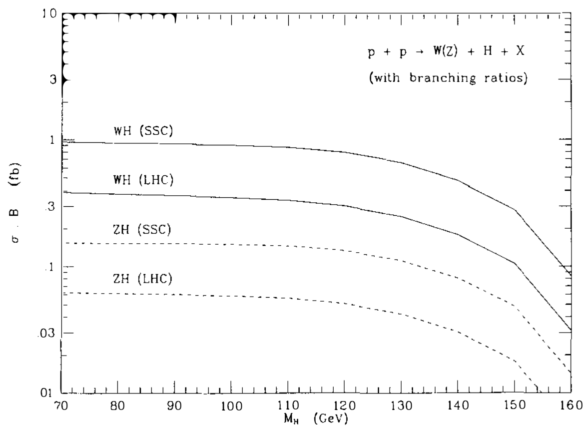

Fig. 10 The signal rate for pp $\to$ (H $\to \gamma\gamma$)W($\ell\nu$)X or pp $\to$ (H $\to \gamma\gamma$)Z($\ell\bar{\ell}$)X at the LHC and SSC (from Ref. [30]).

A similar process, studied in Refs. [35], is pp $\to t\bar{t}H$ followed by H $\to \gamma\gamma$ and the leptonic decay of at least one W from the t, $\bar{t}$ disintegrations. The signal at the LHC is larger by a factor of about 2 with respect to the WH channel, while the background is the same.

The possibility of detecting the Higgs via H $\to \tau^+\tau^-$ as proposed in Ref. [41] was also considered in detail. The conclusion is negative: this channel turns out to be hopeless for the standard Higgs [30]. As we shall see in Section 4 for particular values of the parameters, it could be of use for the SUSY Higgs A.

Turning to the case of a heavy Higgs, $m_H > 2m_Z$, the golden channel is H $\to$ ZZ $\to 4\ell^\pm$ [37], while H $\to$ WW $\to \ell\ell\nu\nu$ is much more difficult, particularly because of the $t\bar{t} \to WWb\bar{b}$ background. The rate for H $\to$ ZZ $\to 4\ell^\pm$ is displayed in Fig. 11 as a function of m_H and m_t [30]. Detailed studies and simulations of the irreducible background from $q\bar{q}$, gg $\to$ ZZ (which is the dominant one in this case) and of the reducible background from $t\bar{t}$, $Zb\bar{b}$, and Z + jets were performed [30]. The reducible background is in all cases small after cuts. With $\int L\ dt \simeq 10^5$ pb^{-1} and $\ell = e, \mu$, the discovery range at the LHC extends up to m_H = 800 GeV (Fig. 12) (with $\int L\ dt \simeq 10^4$ pb^{-1} the corresponding value would go down to 400 GeV). The ultimate discovery range at the LHC could be improved, perhaps, up to $m_H \simeq 1$ TeV by using H $\to$ ZZ $\to \ell\ell\nu\nu$, but the possibility of extracting the signal from the background from $b\bar{b}$, $Zb\bar{b}$, etc., is not demonstrated. Alternatively one could try to use H $\to$ WW $\to \ell\nu$jj or H $\to$ ZZ $\to \ell\ell$jj with jet tagging. Jet tagging was first studied in Ref. [38] and further considered at the Aachen Workshop by M. Seymour [30]. At large m_H, a substantial fraction of the Higgs events is produced by WW fusion. As is well known, the idea of tagging is to detect the near forward and backward quark jets, with $E_\ell \sim O(1$ TeV) and $p_T \sim O(m_W)$ left out after W emission. Studies done at the Workshop indicate that jet tagging may indeed be possible, perhaps even at $L > 10^{33}$ cm^{-2} s^{-1}.

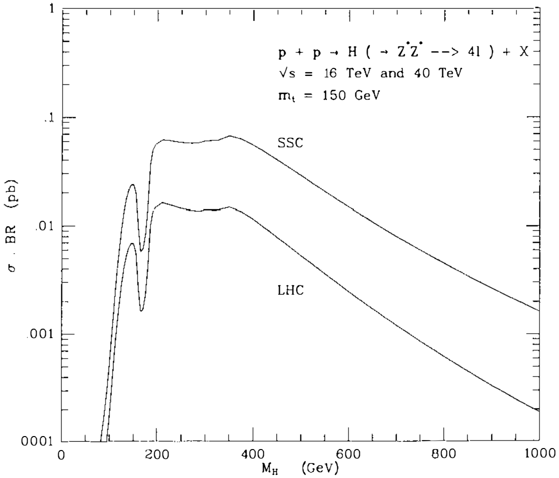

Fig. 11 Signal for H $\to$ ZZ $\to 4\ell^\pm$ at the LHC and SSC (from Ref. [30]).

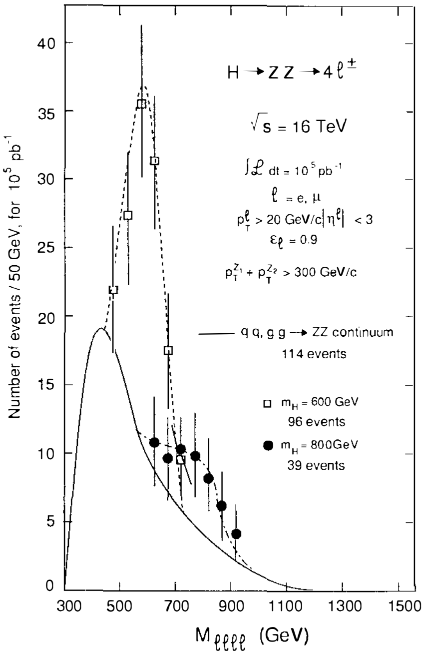

Fig. 12 Signal versus background for $\mathrm{H} \to \mathrm{ZZ} \to 4\ell^{\pm}$ for $m_{\mathrm{H}} = 0.6$–0.8 TeV (from Ref. [30]).

In conclusion, as was stated by D. Froidevaux in Ref. [30], at the LHC, with 10^5 pb^{-1}, the process $\mathrm{H} \to \mathrm{ZZ} \to 4\ell^{\pm}$ with real or virtual Z, allows the range $m_{\mathrm{H}} = 130$–800 GeV to be covered. The same range is obtained at the SSC with 10^4 pb^{-1}. For $m_{\mathrm{H}} = 80$–130 GeV the channels $\mathrm{H} \to \gamma\gamma$ $\mathrm{HW} \to \gamma\gamma\ell\nu$ and $\mathrm{t\bar{t}H} \to \ell\nu\gamma\gamma$ are extremely difficult but feasible. The ratio $\mathrm{S}/\sqrt{\mathrm{B}}$ is actually better at the LHC with 10^5 pb^{-1} than at the SSC with 10^4 pb^{-1}, but the operation at a luminosity 10 times larger is more demanding for the detector. The ultimate discovery range at the LHC could perhaps be extended up to 1 TeV by using $\mathrm{H} \to \mathrm{ZZ} \to \ell\ell\nu\nu$ or $\mathrm{H} \to \mathrm{WW} \to \ell\nu\mathrm{jj}$ with jet tagging, but this is not established.

3 Longitudinal $\mathrm{W}^{\pm}$ and Z

The V states ($\mathrm{V} \equiv \mathrm{W}^{\pm}, \mathrm{Z}$) with helicity zero (longitudinal V, denoted by $\mathrm{V_L}$) are absent in the symmetric limit where the V are massless. It is thus clear that the longitudinal modes are directly related to the symmetry-breaking mechanism. If the Higgs is not found in the LHC discovery range, then the VV interactions become strong and the perturbative cross-section violates unitarity [8] for $m_{\mathrm{H}}, \sqrt{s} \geq \mathrm{O}(1\ \mathrm{TeV})$. This is due to the growth of the $\mathrm{V_L V_L} \to \mathrm{V_L V_L}$ scattering amplitudes, which become dominant in that regime. If the Higgs is not found at the

LHC, the study of the interactions among V_L becomes the most direct way of attacking the symmetry-breaking problem [39]. In a theory with spontaneous symmetry breaking, no matter if the breaking is dynamical (e.g. due to condensates) or induced by either elementary or composite Higgses, the longitudinal V arise from the Goldstone bosons with the corresponding quantum numbers. In fact, at large energies, when contributions of order $m_V/\sqrt{\hat{s}}$, arising from mass terms, can be neglected, the amplitudes for V_LV_L scattering approach those for the corresponding Goldstone bosons ($\sqrt{\hat{s}}$ being the V_LV_L centre-of-mass energy). For example

$$A(W_L^{\pm}Z_L \to W_L^{\pm}Z_L) = A(w^{\pm}z \to w^{\pm}z) + O\left(\frac{m_V}{\sqrt{\hat{s}}}\right) , \tag{3.1}$$

where v_i is the Goldstone boson which corresponds to V_{iL}. This 'equivalence theorem' [40], valid to all orders of perturbation theory, is also used as a handy method for practical computations.

At low momenta, the Goldstone boson couplings are fixed by the symmetry. As a consequence, there are low-energy theorems [41] that specify the Goldstone boson amplitudes at threshold. An effective Lagrangian formalism can be based on the low-energy theorems. This provides a framework for an extrapolation near threshold of the amplitudes which satisfy the low-energy theorems. At $\sqrt{\hat{s}} \gg m_V$ but not too large, one may think to combine the equivalence theorem and the low-energy limit and to apply the effective Lagrangian results directly to V_LV_L scattering. Such smooth extrapolations can provide reasonable approximations only for $\sqrt{\hat{s}} \ll 4\pi G_F^{-1/2}$, provided that no resonances are met on the way. For example, in the Standard Model the regime of low-energy theorems is no longer valid for $\sqrt{\hat{s}} \simeq m_H$, because m_H is a resonance in the VV channel. At large $\sqrt{\hat{s}}$ ($\sqrt{\hat{s}} \gg m_H$), the Higgs contribution cancels [8] the bad high-energy behaviour—obtained by extrapolating the trend derived from the low-energy theorems—which eventually would violate unitarity. For a light Higgs, the high-energy V_LV_L scattering amplitudes remain small, of order $G_F m_H^2$ [8]. In the absence of the Higgs, some other mechanism, which one would like to discover, should intervene to quench the singular high-energy behaviour.

An analogy with QCD can be established: $W_L^{\pm}$ and Z_L are analogous to $\pi^{\pm}$ and π^0 in QCD because V_L are eaten up Goldstone bosons of SU(2) $\otimes$ U(1), while the pions are the (pseudo)-Goldstone bosons of SU(2) $\otimes$ SU(2) chiral symmetry. The pions obey Weinberg's low-energy theorems [42], which are embodied in the formalism of chiral Lagrangians [43]. The chiral Lagrangian regime would hold up to $\sqrt{\hat{s}} \ll 4\pi F_\pi \simeq 1.2$ GeV were it not for the presence of vector mesons ρ, ω that induce drastic differences already at $\sqrt{\hat{s}} \sim m_\rho$. In the case of $W_L^{\pm}$ and Z_L, F_π is replaced by $G_F^{-1/2}$.

Two broad possibilities emerge and have been amply discussed in the literature. On the one hand, the situation of the Standard Model can be stretched up to large m_H, where a very broad enhancement is present in the scalar channel with $I = 0$ (I: weak isospin). On the other hand, the QCD picture can be mimicked with vector resonances with $I = 1$ (ρ) or $I = 0$ (ω). This is for example the case of models based on SU(N_{TC}) technicolour [18–20] or scaled-up QCD (i.e. $N_{TC} = 3$) or the 'BESS' model of Casalbuoni, Gatto et al. [44], which is a non-renormalizable

Lagrangian model with no Higgs (eliminated as in the non-linear σ model [45]) extended to include an extra SU(2), which leads to heavy vector ρ-like states (with $I = 1$).

A more general approach that can generate a QCD-like or a Higgs-like model, or other cases as well, was adopted by Dobado, Herrero and Terron [46] (for related work, see also Ref. [47]). Higher-order terms in the momenta are added to the lowest-order effective Lagrangian. While the lowest-order effective Lagrangian is fixed by the low-energy theorems in terms of a single-energy parameter, $G_F^{-1/2}$, the next-order couplings depend on two arbitrary parameters. By varying those constants one can switch from one type of physics to another. Some procedure of unitarization is implemented in order to extend the model at large $\sqrt{s}$ (in a purely phenomenological way) so that the model formally makes sense also in the presence of resonances.

Extensive studies based on the various models listed above were performed for the Aachen Workshop [30], together with detailed experimental simulations, in order to evaluate the capabilities of the LHC in this domain of physics. The general procedure is to compute the VV scattering amplitudes in a given model, to compare the results with the Standard Model prediction for some large but still admissible Higgs mass, to check whether the deviations would be measured at the LHC, and to disentangle the different models.

The processes that are best suited for an experimental investigation are those with no $t\bar{t} \to WWb\bar{b}$ background: ZZ, $W^{\pm}Z$, and $W^{\pm}W^{\pm}$ (equal charges!) final states. Different qualitative behaviours are expected in these channels, depending on the dynamics of V_LV_L scattering: in the Higgs-like regime, sizeable effects are expected in the ZZ channel and not in the $W^{\pm}Z$ or $W^{\pm}W^{\pm}$ reactions. Conversely a ρ-like resonance would show up in the WZ channel and not elsewhere.

For equal-sign WW final states, the production rate of W^-W^-, with $M_{WW} > 0.8$ TeV, is about one third of that of W^+W^+ (because u quarks are more abundant than d quarks at large x in the proton). The background from $W^{\pm}t\bar{t}$, from $q\bar{q} \to W^{\pm}W^{\pm}q\bar{q}$ via gluon exchange, and from QCD jets has been evaluated by Barger et al. [48]. In models with no $I = 2$ resonances, as those studied by Dobado et al. [46], there is little activity in the channel $pp \to W^{\pm}W^{\pm}X \to \ell^{\pm}\nu\ell^{\pm}\nu X$, and the signal is small with respect to the background (Fig. 13). For $\int L\ dt \simeq 10^5$ pb^{-1} the rate is of the order of 10 events per year at $M_{WW} > 0.8$ TeV. The situation is no better at the SSC with 10^4 pb^{-1}. This does not necessarily mean that this process is not interesting, because the actual dynamics could be different from that of the models studied at the Workshop. If a doubly-charged resonance exists, it would show up in this channel.

In the $ZZ \to 4\ell^{\pm}$, $\ell = e, \mu$, channel the signal from $W_LW_L \to Z_LZ_L$ plus $Z_LZ_L \to Z_LZ_L$ was computed in the model by Dobado et al. [46], compared with the irreducible background from the Standard Model processes $q\bar{q}$, $gg \to ZZ$. For $\int L\ dt \simeq 10^5$ pb^{-1}, $M_{ZZ} > 0.5$ TeV, $p_T^Z >$ 10 GeV, $|y_Z| < 2.5$, the background amounts to about 220 events (for $m_t = 100$ GeV), while the signal is of about 15 events in a Higgs-like picture, and half of that in a scaled-up QCD model. The corresponding numbers at the SSC, with 10^4 pb^{-1} and the same cuts, are 73 (background), 10 (Higgs-like), and 5 (QCD-like) events. Without jet tagging it is difficult

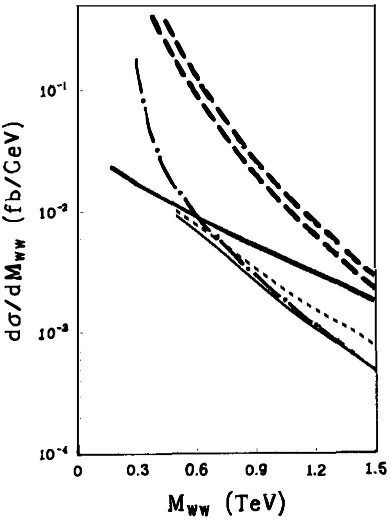

Fig. 13 Like-sign WW invariant-mass distribution for various strongly-interacting models and for the total background. Rates for W^+W^+ and W^+W^- are added [46]. The short-dashed line is for the QCD-rescaled case and the lower solid line is for the Higgs-like model. The upper solid curve corresponds to the unitarized-LET results [48]. The Standard Model rates for $M_H = 1$ TeV are also displayed for comparison [48] (dot-dashed line). The long-dashed lines are the predictions for the total background in the cases $m_t = 100$ GeV (upper line) and $m_t = 200$ GeV (lower line), respectively [48].

to separate the VV $\to$ VV signal from the irreducible background, particularly because the latter is only computable with limited accuracy and both the signal and the background have a structureless mass distribution. (Recently, the next-to-leading QCD corrections to $q\bar{q} \to ZZ$ have been computed [49].)

The prospects are much more promising for models with resonances, as for example a $\rho^\pm$-like particle observable in $W^\pm Z$ final states or an ω-like object visible in $Z\gamma$. At the Aachen Workshop the WZ channel was studied in full detail in SU(N_{TC}) models realized in the effective Lagrangian approach [46] and in the BESS model [44, 50]. In scaled-up QCD

$$m_{\rho_{TC}} \simeq \frac{v}{F_\pi} m_\rho \sim 2 \text{ TeV} \tag{3.2}$$

$$\Gamma(\rho_{TC} \to VV) \simeq \frac{m^3_{\rho_{TC}}}{96\pi v^2} \simeq 450 \text{ GeV}. \tag{3.3}$$

For SU(N_{TC}) with $N_{TC} \neq 3$ one takes [18]:

$$m_{\rho_{TC}} \simeq 2 \text{ TeV}\sqrt{\frac{3}{N_{TC}}}, \quad \Gamma(\rho_{TC} \to VV) \simeq 450 \text{ GeV}\left(\frac{3}{N_{TC}}\right)^{3/2}. \tag{3.4}$$

Thus for $N_{TC} \simeq 12$ one has $m_{\rho_{TC}} \simeq 1$ TeV and $\Gamma_{\rho_{TC}} \simeq 55$ GeV, while for $N_{TC} = 5$, $m_{\rho_{TC}} \simeq 1.5$ TeV and $\Gamma_{\rho_{TC}} \simeq 185$ GeV. The results for these three representative cases ($m_{\rho_{TC}} = 1, 1.5, 2$ TeV) are summarized in Fig. 14. The full process under consideration is $pp \to W^\pm ZX \to$

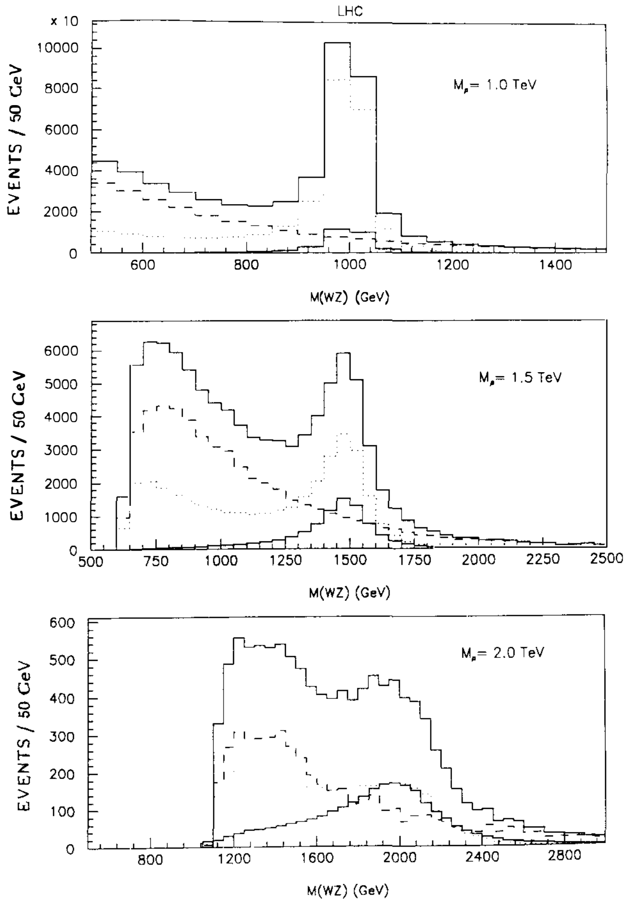

Fig. 14 WZ invariant-mass distribution for the signal and background processes with the optimal cuts (a 2.5 rapidity cut has been chosen) [46]. Rates are for W^+Z + W^-Z and for $L = 4 \times 10^5$ pb^{-1}. The results of the signal are for three possible cases in SU(N_{TC}) theories corresponding to: $m_\rho = 1.0$ TeV, 1.5 TeV, and 2.0 TeV, respectively. The lower solid histogram represents the WZ fusion contribution to the signal. The dotted histogram is the $q\bar{q}'$ annihilation contribution to the signal via ρ–ω mixing. The total background is the dashed histogram, and the total signal + background is the upper solid histogram.

$\ell^\pm \nu \ell^+ \ell^- X$, $\ell = e, \mu$. The W^+Z rate is about twice the W^-Z rate. The resonant signal in $W_L^\pm Z_L$ is produced either by WZ fusion or by $q\bar{q}$ annihilation with ρ_{TC} coupled via W–ρ_{TC} mixing. The irreducible background is from the standard processes $q\bar{q} \rightarrow WZ$, $W\gamma \rightarrow WZ$, and $WZ \rightarrow WZ$ (with no ρ_{TC} exchange). With optimized cuts the following S/B ratios were obtained at the LHC [46] (in number of events per 10^5 pb^{-1}): 660/53 for $m_{\rho_{TC}} = 1$ TeV, 50/11 for $m_{\rho_{TC}} = 1.5$ TeV, and 20/13 for $m_{\rho_{TC}} = 2$ TeV. At the SSC with 10^4 pb^{-1} (with different cuts optimized to the SSC case) the corresponding numbers are: 263/24, 36/8, and 24/16, respectively. The resonance is visible in the mass and p_T distributions. The invariant mass distributions in the

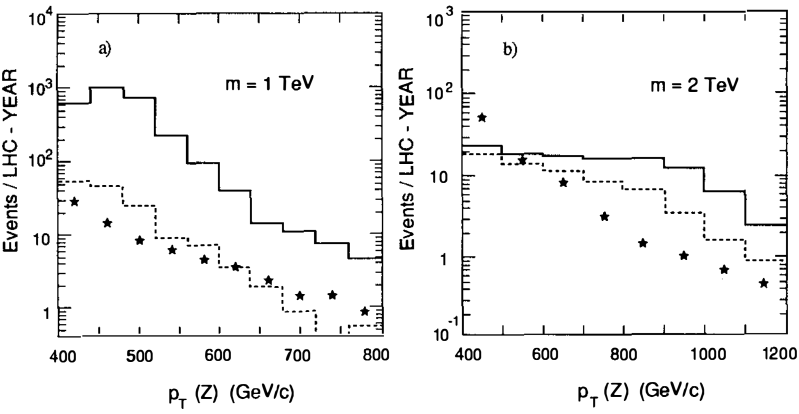

Fig. 15 Signal and background for $\rho^{\pm}_{\rm TC}$, with $m_{\rho^{\pm}_{\rm TC}} \simeq$ 1–2 TeV. The solid line is the total signal in $\rm W^{\pm}_L Z_L$ final states (obtained by adding boson fusion and $\rm q\bar{q}'$ annihilation). The dashed line is the fusion contribution alone. The stars indicate the total background. (From Ref. [30].)

LHC case are shown in Fig. 14. Detailed simulations presented by Rodrigo et al. [30] show that the signal clearly emerges at large $p_{\rm T}$ over the complete background, also including the reducible one with 3-lepton events from $\rm t\bar{t}$ production (Fig. 15).

The production of a different type of ρ-like resonance in the WZ channel was studied in the context of the BESS model [44, 50]. The values of the free parameters $m_{\rho_{\rm TC}}$, g'' were chosen in such a way as to have $m_{\rho_{\rm TC}} = 1$, 1.5, 2, 2.5 TeV with widths 11–44, 84–355, 353, 455 GeV, respectively. The $\rho^{\pm}_{\rm TC}$ is coupled to WZ and also to $\rm q\bar{q}$ via the $\rm W^+$–$\rho^+_{\rm TC}$ mixing (an additional direct coupling with quarks could be switched on by letting a parameter called b be different from zero). The background is the same as in the previous discussion. In Fig. 16a,b, we report the $p_{\rm TC}$ distributions for the case $\rho_{\rm TC} = 2.0$ TeV, with $b = 0$. Here again the LHC with $\int L\ dt \simeq 10^5$ pb^{-1} is compared with the SSC with $\int L\ dt \simeq 10^4$ pb^{-1}. Even if $m_{\rho_{\rm TC}} \simeq 2$ TeV is large enough to provide a special advantage to the SSC, we see that the S/B ratios in the two cases are comparable [684/310 (LHC) and 1010/462 (SSC)]. The discovery range at both LHC and SSC extends up to $\sim$ 2.5 TeV.

Summarizing (see the report by M. Lindner, S. Dimopoulos et al. in Ref. [30]): $\rm W^{\pm}W^{\pm}$ is small, below the background, in models with no $I = 2$ resonances. The ZZ channel is in principle good for the Higgs-like case, but it is very difficult to disentangle a non-resonant signal from the continuum. The $\rm W^{\pm}Z$ (or $\rm Z\gamma$) is good for ρ-like (ω-like) resonances. A $\rho_{\rm TC}$ resonance with $m_{\rho_{\rm TC}} < 2.5$ TeV can be detected in the WZ channel at the LHC with $\int L\ dt \simeq 10^5$ pb^{-1} or at the SSC with $\int L\ dt \simeq 10^4$ pb^{-1}. In conclusion if there are resonances with $\Gamma \ll M$, they can be detected. Otherwise a structureless signal is difficult to be established both at the LHC and at the SSC.

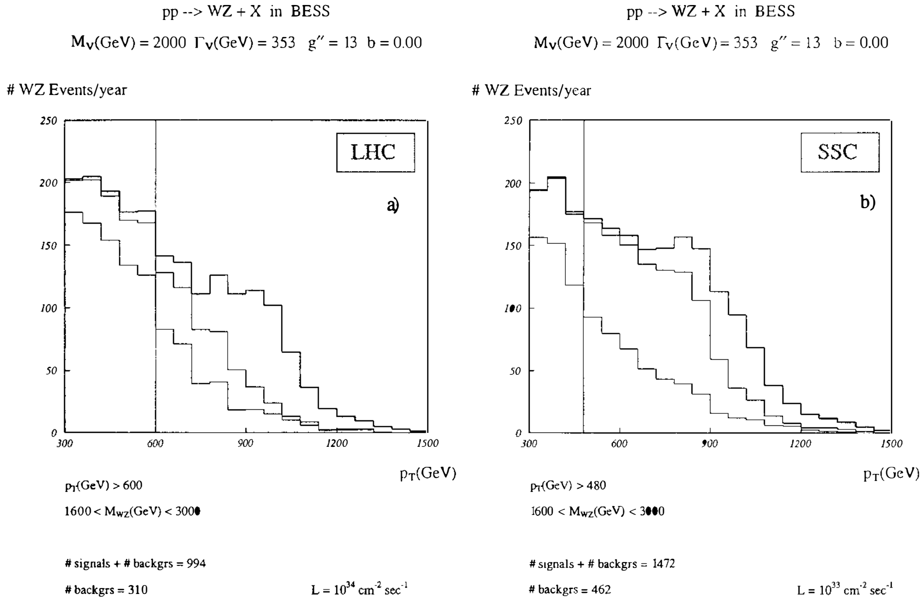

Fig. 16 Transverse-momentum distribution of signal and background at the LHC (a) and the SSC (b), for ρ_{TC} production in BESS (from Refs. [44] and [50]), with $m_{\rho_{TC}} = 2$ TeV and $b = 0$. The upper histogram is the signal from $q\bar{q}$ annihilation, the centre one that from fusion, and the lower one from the total background.

4 Supersymmetric Higgses

As we discussed in the introduction, many theorists consider that fundamental scalar Higgses are most likely to be accompanied by supersymmetry in order to make the theory natural when looked down from very high energy scales such as $M_{\rm GUT}$ or $M_{\rm Pl}$. However in all supersymmetric extensions of the Standard Model at least two Higgs doublets are necessary [17, 29], giving their masses one to the up fermions and the other to the down ones. Thus in supersymmetric models there are at least three neutral and two charge-conjugated charged physical Higgses. In the MSSM as discussed by F. Zwirner [51] the spectrum of physical Higgses is specified by two parameters: the mass of one of the neutral Higgses and tg $\beta = v_{\rm u}/v_{\rm d}$, the ratio of the vacuum expectation values of the Higgses that give mass to up fermions, $v_{\rm u}$, and to down fermions, $v_{\rm d}$. In the MSSM, tg β is always larger than 1 (while in a generic two-doublet model there is no such restriction). Also, values of tg $\beta > m_{\rm t}/m_{\rm b}$ are not allowed. The neutral Higgses are denoted by h, A, and H: h is the lightest Higgs ($J^{\rm CP} = 0^+$), A is the Higgs with opposite CP (0^-), and H is the heavy Higgs with quantum numbers 0^+. At tree level, in terms of the parameters tg β and $m_{\rm A}$, one has [17, 29, 51]

$$m^2_{\rm H^\pm} = m^2_{\rm A} + m^2_{\rm W} \tag{4.1}$$

$$m^2_{\mathrm{h,H}} = \frac{1}{2}\left[(m^2_{\mathrm{A}} + m^2_{\mathrm{Z}}) \pm \sqrt{(m^2_{\mathrm{A}} + m^2_{\mathrm{Z}})^2 - (2m_{\mathrm{A}} m_{\mathrm{Z}} \cos\ 2\beta)^2}\right], \tag{4.2}$$

so that $m_{\mathrm{H}^\pm} \geq m_{\mathrm{W}}; m_{\mathrm{h}} \leq m_{\mathrm{Z}}, m_{\mathrm{A}}; m_{\mathrm{H}} > m_{\mathrm{Z}}, m_{\mathrm{A}}$. From LEP data analysed in terms of tree level formulae one would obtain tg $\beta \geq 1.6, m_{\mathrm{A}} \geq 40$ GeV, $m_{\mathrm{h}} \geq 33$ GeV. But for large m_{t} there is the possibility that radiative corrections could induce rather large shifts in the Higgs masses [52]. In particular m_{h} could exceed m_{Z}. The results of Refs. [52] imply that for $m_{\mathrm{t}} \sim$ 130 GeV, the shift of m_{h} due to the radiative corrections is of a few GeV and becomes rapidly larger with increasing m_{t} (the effect increases as m_{t}^4).

The case of the MSSM is a particularly important and interesting one. The implications on the LHC of the possibility that the supersymmetric version of the Higgs sector is realized in nature have been discussed at the Aachen Workshop by Kunszt and Zwirner [30]. However their paper in Ref. [30] is affected by some numerical errors and their conclusions have been changed [53]. The first observation is that if the MSSM is true, then most probably a Higgs will be found at LEP. We repeat that the lucky event of the discovery of a Higgs particle at LEP does not in any sense diminish the physics case for the LHC. This is obviously true if the observed properties of the light Higgs depart from the behaviour of the standard Higgs and are consistent with the MSSM. But this is also true if the accessible information obtained at LEP on the light Higgs is compatible with the Standard Model. In fact for most of the parameter space the properties of the light Higgs are close enough to the Standard Model for LEP not to be able to clearly distinguish the two cases. It is only the experimental investigation of the LHC energy domain that can possibly clarify the issue. In particular the question of the search for the SUSY Higgses at the LHC is an important one.

The production cross-sections of SUSY Higgses at the LHC/SSC are often larger than for the Standard Model Higgs of the same mass. This is because of the addition of s-quark loops in the gluon-fusion mechanism and also because of the larger couplings to $\mathrm{b\bar{b}}$ for tg β large. However the couplings to $\gamma\gamma$, WW, and ZZ are typically suppressed with respect to the Standard Higgs. For example the modes A $\to$ WW or ZZ are forbidden for the A boson. But, for large tg β, the channels A $\to \tau\tau$ or H $\to \tau\tau$ are promising: the $\tau\tau$ mode which is no good for the Standard Higgs can be viable for SUSY Higgses. Similarly the mode H $\to$ ZZ $\to 4\ell^\pm$ is good for the heavy SUSY Higgs provided that $m_{\mathrm{H}} < 2m_{\mathrm{t}}$ and tg β is small.

In conclusion, for the neutral Higgses of the MSSM the detection is in general a hard problem. A separate analysis, as complicated as the one for the Standard Higgs, would be necessary for each set of values of tg $\beta, m_{\mathrm{A,H}}$, and m_{t}. While much more work is needed on this subject, there certainly are windows in the parameter space where detection is possible for at least either A or H.

The case of the charged Higgs was also considered at the Aachen Workshop, especially by M. Felcini [30]. The charged Higgs could be observed if it is present in t decays: t $\to \mathrm{H^+b}$. As the dominant $\mathrm{H}^\pm$ decay would be $\mathrm{H}^+ \to \tau^+\nu$, the signature would be a measurable violation of τ–μ universality in $\mathrm{t\bar{t}}$ events. For $m_{\mathrm{t}} \simeq 200$ GeV, m_{H^+} could be detected in the range $m_{\mathrm{H}^\pm} =$ 100–150 GeV, while for $m_{\mathrm{t}} = 150$ GeV, m_{H^+} would be visible up to $m_{\mathrm{H}} \sim 100$ GeV (Fig. 17).

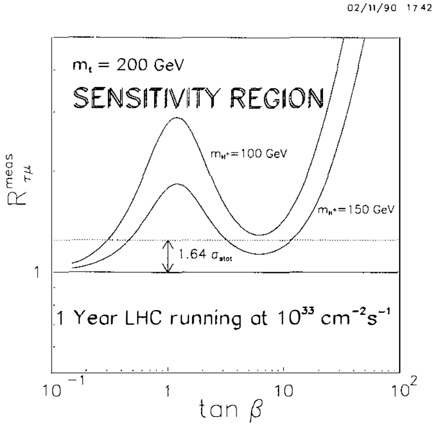

Fig. 17 Sensitivity to violations of the $\tau \leftrightarrow \mu$ universality induced by a charged Higgs (from M. Felcini).

5 Conclusion

The main goal of experiments on particle physics in the near future is the clarification of the electroweak symmetry-breaking problem. The solution must be within the TeV energy region: the origin of the weak scale cannot lie too far from $G_F^{-1/2} \sim 293$ GeV. Probably a whole universe of new physics will open up. Examples are offered by the supersymmetric model, which provides a well-defined extension of the Standard Model, which is more natural than the Standard Model itself. Other possibilities are less well defined. Apart from possible completions of the Standard Model in the direction of extending the electroweak group (new W′ and Z′), all alternatives to fundamental scalar Higgses and supersymmetry involve new strong forces and a breakdown of the perturbative regime in the TeV energy region.

A common feature of all conceivable ways beyond the Standard Model is the prediction of a rich spectroscopy of new states and new phenomena. This means that one expects discoveries over a wide range of energies. Actually it would be a great thing for the LHC and SSC if the low-lying fringes of the new spectroscopy were already found at LEP 1 and 200. Far from decreasing the physics motivations in favour of the LHC and SSC, the discovery at LEP of some new physics or at least of some departures from the Standard Model would make the argument for the LHC even stronger.

The results obtained at the Aachen Workshop clearly demonstrate that the discovery potential of the LHC with $L \simeq 10^{34}$ cm^{-2} s^{-1} is perfectly adequate to the goal of solving the problems of the electroweak symmetry breaking and of the origin of the weak scale of mass. It is also evident, from the detailed comparison made in the previous sections, that the LHC with $L = 10^{34}$ cm^{-2} s^{-1} is very much comparable with the SSC with $L = 10^{33}$ cm^{-2} s^{-1}. For standard Higgses, we have seen that the discovery range extends up to (0.8–1) TeV in both cases. In WW, WZ, ZZ scattering resonances, such as the ρ-like or ω-like ones, vector bosons of technicolour are visible up to 2–2.5 TeV at the LHC and the SSC, while in both cases non-resonant amplitudes are very difficult to study. New W' and Z' can be found up to 4.5–5 TeV at the LHC, and up to 5–6 TeV at the SSC. Gluinos and squarks can be observed up to 1.5 TeV at the LHC and up to 1.5–2 TeV at the SSC.

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