Title
THE QUEST FOR THE KEGGS BOSON

Author
Hinchliffe, I.

Publication Date
1987-02-01

LBL-23038 c.2

Lawrence Berkeley Laboratory
UNIVERSITY OF CALIFORNIA

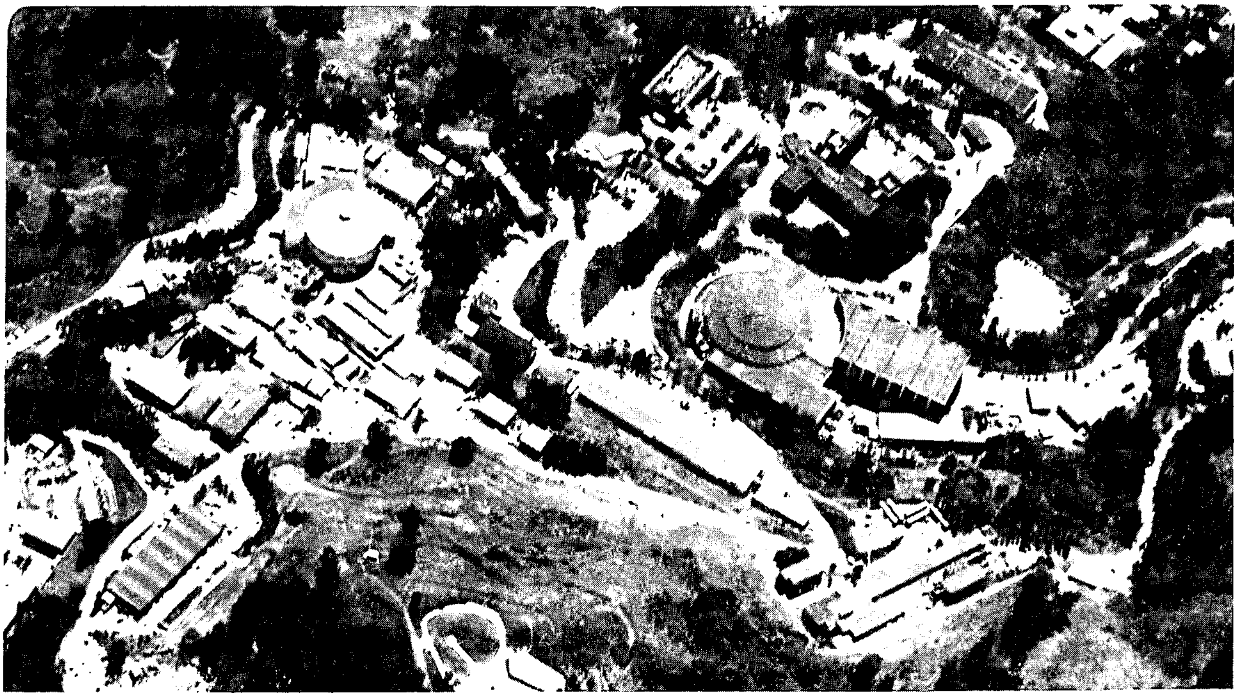
Physics Division

Invited talk presented at the 2nd Annual
Lake Louise Winter Institute on Electroweak Interactions,
Lake Louise, Canada, February 15–21, 1987

THE QUEST FOR THE HIGGS BOSON

I. Hinchliffe

February 1987



LBL-23038 c.2

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

February 1987

LBL-23038

THE QUEST FOR THE HIGGS BOSON*

I. Hinchliffe
Lawrence Berkeley Laboratory
University of California
Berkeley, California 94720

Invited talk at the 2nd Annual Lake Louise Winter Institute, Feb. 15–21 1985.

---

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

Abstract

A review is given of the various ways which have been suggested to search for the Higgs boson of the Weinberg–Salam model.

There is very little experimental information about the Higgs sector of the $SU(2) \times U(1)$ model[1] other than that it must have a custodial $SU(2)$ symmetry so that $\rho = \frac{M_W}{M_z \cos\theta_W} = 1$ at tree level. Do we know anything from theoretical studies?

Since $m_H^2 = \frac{2\lambda}{g_2^2} M_W^2$, it would appear that m_H could be made arbitrarily small by reducing the Higgs self coupling λ. This is not the case since for very small λ one must consider the effect of gauge interactions which induce Higgs self interactions at higher order. The Higgs self interactions are described by the effective potential

$$V_{eff}(\Phi) = -\mu^2 \Phi^+ \Phi + \lambda(\Phi^+ \Phi)^2. \tag{1}$$

Radiative corrections from Feynman diagrams of the type indicated in Fig. 1 modify this potential. At one loop $V_{eff}(\Phi)$ becomes[2]

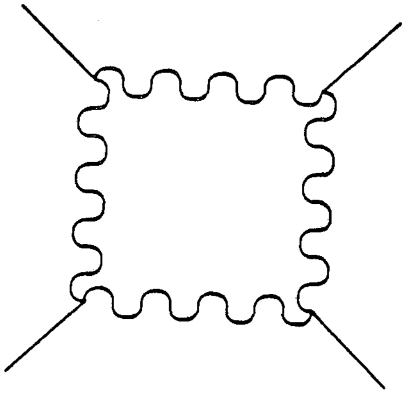


Figure 1: Feynman diagram showing a contribution to the effective potential for the Higgs field due to interactions of the Higgs with W bosons.

$$V(\Phi) = -\mu^2(\Phi^+\Phi) + \lambda(\Phi^+\Phi)^2 + c(\Phi^+\Phi)^2 \log((\Phi^+\Phi)/M^2) \tag{2}$$

where

$$c = \frac{1}{16\pi^2 v^4}(3(2M_W^4 + M_Z^4) + m_H^4 - 4\sum_f m_f^4)$$

where V is the vacuum expectation value of the Higgs field. M is a renormalization scale and m_f is a fermion mass. The Higgs mass m_H is given by

$$m_H^2 = \frac{\partial^2 V}{\partial \Phi^2}\Big|_{\Phi=\langle\Phi\rangle} . \tag{3}$$

In general V_{eff} will have more than one minimum. If we require that the minimum with $\langle\Phi\rangle \neq 0$ is lower than that with $\langle\Phi\rangle = 0$ (the phase in which the W boson remains massless), so that this phase will be the true ground state, then a bound on λ, and hence m_H, can be obtained since all the other quantities in Eq. (2) are known. We have[3]

$$m_H \gtrsim 7\ GeV.$$

A more detailed study which requires that the universe not be trapped at $\langle\Phi\rangle = 0$ for too long[4] gives $m_H \gtrsim 10\ GeV$. This bound is extremely model dependent. A similar bound will exist in models with different Higgs sectors.[5] In models with an arbitrary number of Higgs doublets there must be at least one physical Higgs boson with a mass greater than this bound.

As λ is increased m_H increases. Eventually λ will become too large for the perturbative formula for the Higgs mass to be valid. We can estimate this value naively by requiring that $\lambda^2/4\pi$ be less than one. This implies $m_H \lesssim 600\ GeV$. In order to be more precise it is necessary to consider the effects of the constraints imposed by partial wave unitarity.[6]

Consider the S matrix for a two-particle scattering process $a + b \to c + d$. Unitarity requires that

$$S^+S = 1. \tag{4}$$

Writing $S = 1 + iT$, we have

$$-ImT = T^+T. \tag{5}$$

The scattering matrix T is given by

$$T = (2\pi)^4\delta^4(p_a + p_b - p_c - p_d)\frac{1}{(2\pi)^6}\frac{1}{s}\left|M_{ab\to cd}\right|^2 . \tag{6}$$

Here p_i is the momentum of particle i and M_i the invariant matrix element obtained, for example, by calculating a set of Feynman diagrams. M may be decomposed as follows:

$$M(s, \cos\theta) = 16\pi \sum_{J=0}^{\infty} (2J+1) A_J(s) P_J(\cos\theta). \quad (7)$$

θ is the center-of-mass scattering angle between particles a and c, $P_J(\cos\theta)$ is a Legendre polynomial and $A_J(s)$ is some function. Equation 5 implies that

$$Im\ A_0 \geq |A_0|^2. \quad (8)$$

We can expand A_0 as a perturbation series in some coupling constant g

$$A_0(s) = a_1(s) g^2 + a_2(s) g^4 + \ldots \quad (9)$$

If the perturbation expression is reliable then

$$g^2 < a_1/a_2. \quad (10)$$

The Born term $a_1 g^2$ is real, hence Eq. 8 implies that

$$-Im(a_2 g^4) > (a_1 g^2). \quad (11)$$

But $|A| > |ImA|$ for any A, so that the requirement that perturbation theory be reliable implies that

$$\left|a_1 g^2\right| < 1. \quad (12)$$

Now this result can be applied to the process $H + H \rightarrow H + H$. If we assume that $m_H >> M_W$, the relevant Feynman diagrams are shown in Fig. 2 and give

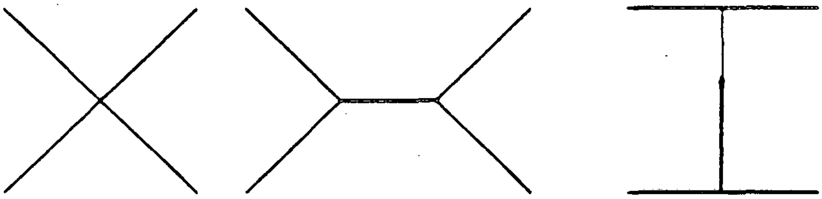


Figure 2: Feynman diagram for the process $HH \rightarrow HH$ which dominate in the limit $m_H >> m_W$.

$$A_0(HH \rightarrow HH) = -\frac{G_F m_H^3}{8\pi\sqrt{2}} \left[3 + \frac{9 m_H^2}{s - m_H^2} - \frac{2 m_H^2}{s - m_H^2} \log(s/m_H^2 - 3) \right]. \quad (13)$$

Requiring $|A_0| < 1$ (see Eq. (12)) in the limit $s \rightarrow \infty$ implies that

$$m_H < 1.7\ TeV.$$

A stronger bound is obtained by considering the coupled channel problem: $HH \to ZZ, HH \to WW, WW \to ZZ, HZ \to HZ, HW \to HW$. In this case one has[6,7]

$$m_H < \sqrt{\frac{8\pi\sqrt{2}}{3G_F}} = .98\ TeV. \tag{14}$$

Arguments based in the triviality of scalar field theory lead to similar conclusions.[8] This bound indicates that there must be a scalar particle of mass less than 1 TeV or the Weinberg-Salam model will contain a strong, non-perturbative coupling. The presence of such a coupling implies that there must be non-perturbative structure in the WW or ZZ channel for WW or ZZ invariant masses of order 1 TeV. (Recall that the longitudinal components of the W and Z come from the Higgs fields.)

The basic argument that I have just outlined contains the essential features which justify the choices of energy and luminosity for the SSC. In order to probe the nature of the interactions responsible for the breakdown of the $SU(2) \times U(1)$ symmetry it is necessary to probe the WW and ZZ system with invariant masses of order 1 TeV.

Let us now turn to the possible experimental signatures for Higgs bosons. The Higgs can decay to fermion anti-fermion (of mass m_f), WW and ZZ final states with the following partial widths:

$$\begin{aligned}
\Gamma(H \to f\bar{f}) &= \frac{G_F m_f^2 m_H}{4\pi\sqrt{2}}(3)(1 - 4m_f^2/m_H^2)^{3/2}, \\
\Gamma(H \to W^+W^-) &= \frac{G_F m_H^3}{32\pi\sqrt{2}}(4 - 4\epsilon + 3\epsilon^2)(1-\epsilon)^{1/2}, \\
\Gamma(H \to ZZ) &= \frac{G_F m_H^3}{64\pi\sqrt{2}}(4 - 4\epsilon' + 3\epsilon'^2)(1-\epsilon')^{1/2}.
\end{aligned} \tag{15}$$

with $\epsilon' = 4M_Z^2/m_H^2$ and $\epsilon = 4M_W^2/m_H^2$. The factor of 3 is included in the first expression only if f is a quark. The implications of these formulae are easy to see. If $m_H < 2M_W$, the Higgs will decay dominantly into the heaviest fermion channel which is open. Once m_H is greater than $2M_W$, the decay into two gauge bosons will dominate. This effect is shown in Fig. 3. Notice that the width grows rapidly as m_H is increased. Eventually $\Gamma/m_H \sim 0(1)$: this is another manifestation of the breakdown of perturbation theory at large values of m_H.

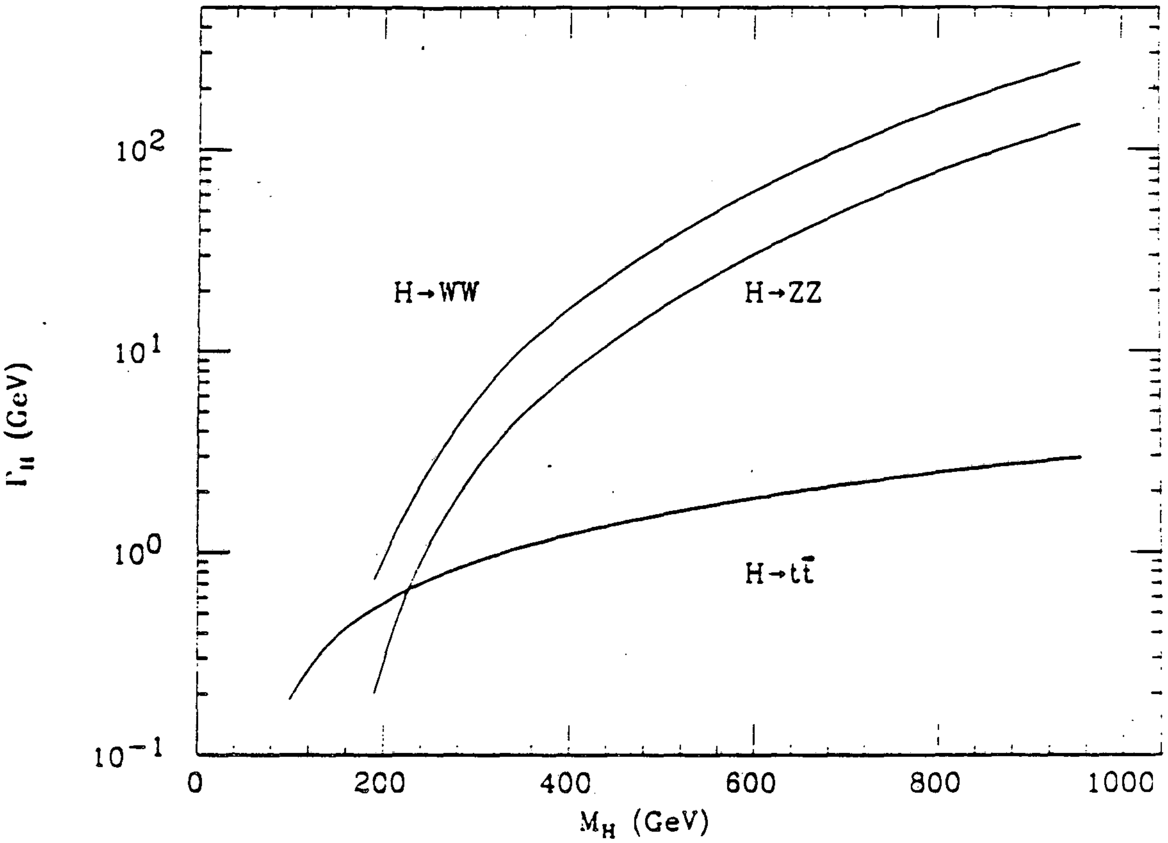


Figure 3: The partial widths $H \to t\bar{t}$ (solid lines), W^+W^- (dashed line) and ZZ (dotted line) as a function of m_H. The top quark mass is taken to be 40 GeV.

The Higgs can be produced in e^+e^- annihilation from the decay of a Z through the graph shown in Fig. 4, with a rate shown in Fig. 5. The rate is given by[9]

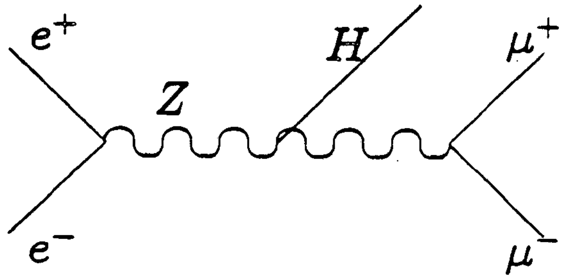


Figure 4: Feynman diagram for the process $e^+e^- \to Z \to H + \mu^+\mu^-$.

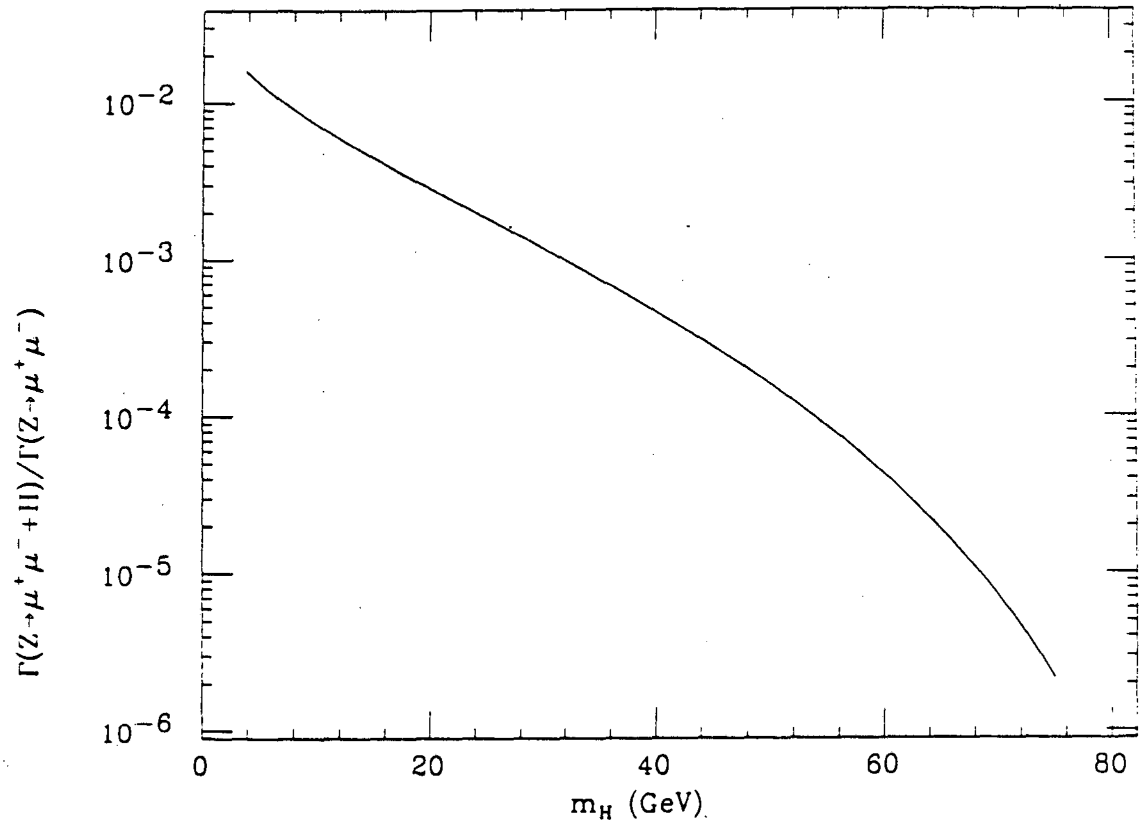


Figure 5: The ratio of widths $\Gamma(Z \to H\mu^+\mu^-)/\Gamma(Z \to \mu^+\mu^-)$ as a function of the Higgs mass.

$$\frac{1}{\Gamma(Z \to \mu^+\mu^-)} \frac{d\Gamma(Z \to H + \mu^+\mu^-)}{dx} = \frac{\alpha}{4\sin^2\theta_W\cos^2\theta_W}$$

$$\times \frac{\left(1 - x + \frac{x^2}{12} + \frac{2m_H^2}{3M_Z^2}\right)\left(x^2 - \frac{4m_H^2}{M_Z^2}\right)^{1/2}}{\left(x - \frac{m_H^2}{M_Z^2}\right)^2} \qquad (16)$$

Here $x = 2E_H/M_Z$ where E_H is the energy of the Higgs boson. The rate is rather small, but the signature is very clean. It is not necessary to reconstruct the Higgs from its decay products; one searches for a peak in the mass recoiling against the lepton pair. Backgrounds arise from the production of a heavy quark pair if both of the quarks decay semileptonically. If we require that the leptons be isolated this background is not important. If one looks in the e^+e^- decay channel, then there is a background from the two photon process $e^+e^- \to e^+e^- +$ hadrons which produces a serious problem for Higgs masses below about 8 GeV. A Higgs of mass less than 40 *GeV* should be discovered at LEP/SLC using this process. If the Higgs mass exceeds 0.6 M_Z, then this rate is exceeded by that from $Z \to H + \gamma$.[10] Again the signal is very clean, but the small rate makes it unlikely that this process will be observed.

The Higgs boson can also be produced at higher energies in e^+e^- annihilation via the process $e^+e^- \longrightarrow Z + H$,[11] with the rate shown in Fig. 6. The production cross-section at center-of-mass energy $\sqrt{s}$ is given by

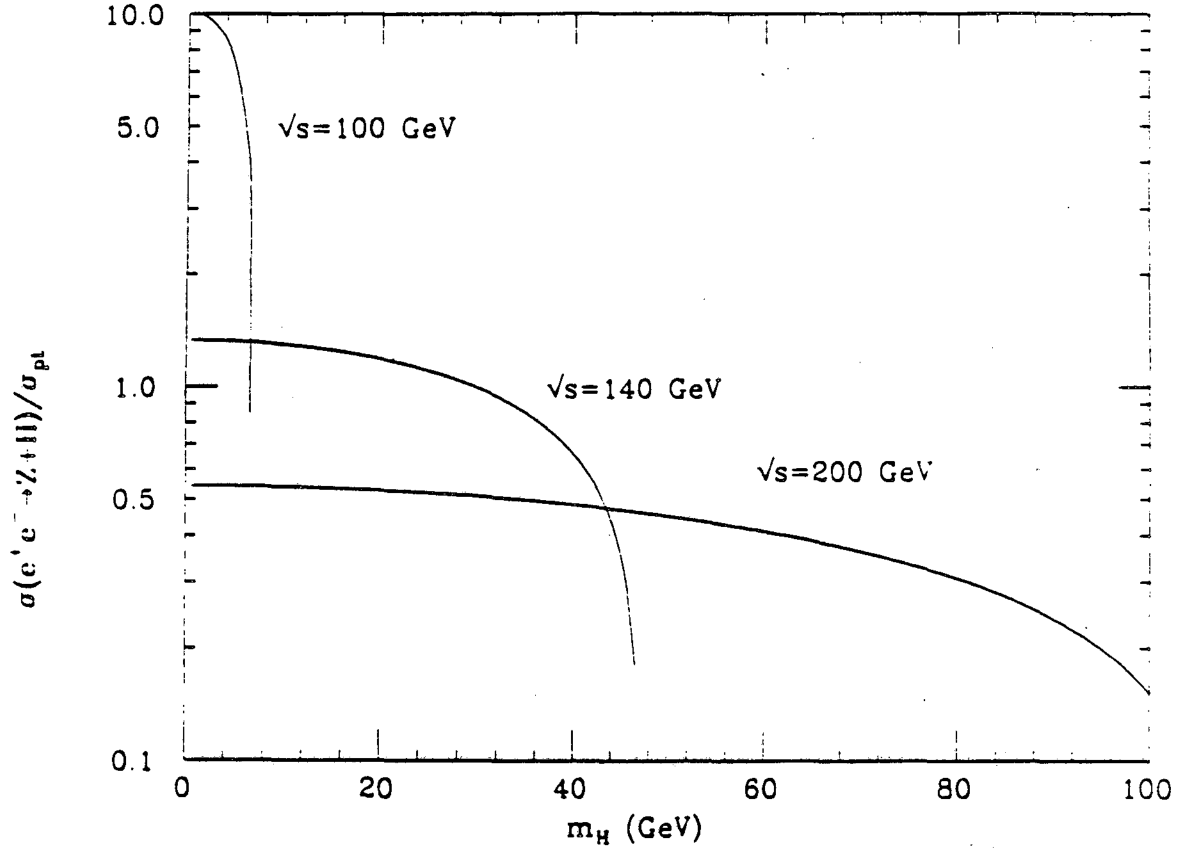


Figure 6: The cross-section for the process $e^+e^- \longrightarrow Z + H$ as a function of m_H for various values of $\sqrt{s}$.

$$\frac{d\sigma(e^+e^- \longrightarrow Z+H)}{d(\cos\theta)} = \frac{\pi\alpha^2(1+8sin^4\theta_W - 4sin^2\theta_W)}{16sin^4\theta_W cos^4\theta_W(s-M_Z^2)^2}\frac{2\kappa}{\sqrt{s}}(M_Z^2 + \frac{\kappa^2 sin^2\theta}{2}). \quad (17)$$

Here θ is the angle between the Higgs and the beam and κ is the Higgs momentum. Again it is not necessary to reconstruct the final state arising from the Higgs decay. The cross section is not large, particularly if the Z can only be detected via its decay to $\mu^+\mu^-$ or e^+e^-. Nevertheless LEP should be able to probe Higgs masses up to $0.9(\sqrt{s} - M_Z)$ using this mechanism.

Another potentially important process is the decay of the toponium bound state (θ) into $H + \gamma$,[12] the rate for which is shown in Fig. 7. Since coupling of a

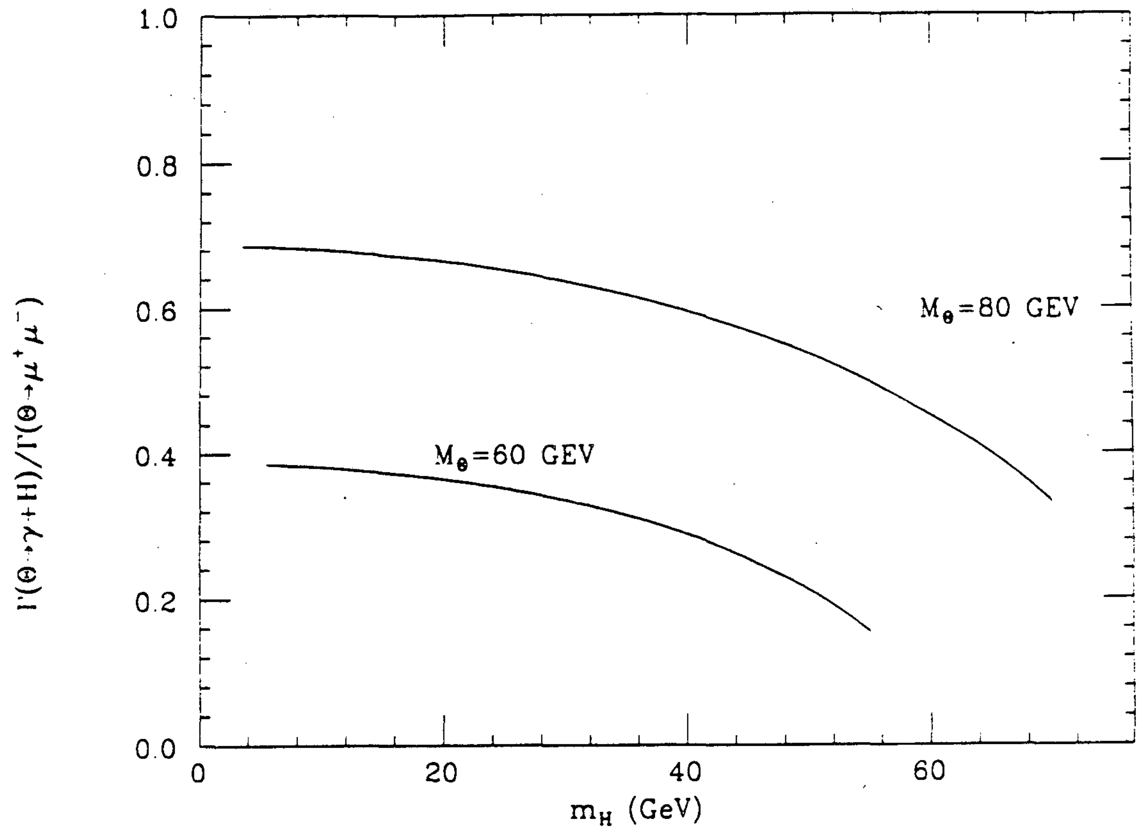


Figure 7: The ratio of decay widths $\Gamma(\theta \rightarrow H\gamma)/\Gamma(\theta \rightarrow \mu^+\mu^-)$ for the decay of the 1^{--} bound state (θ) of $t\bar{t}$. It has been assumed that θ has a mass of 80 *GeV*.

Higgs to a quark is proportional to the quark mass, the rate will be largest if the top quark mass (m_t) is large. The rate is given by

$$\frac{\Gamma(\theta \rightarrow H\gamma)}{\Gamma(\theta \rightarrow \mu^+\mu^-)} = \frac{G_F m_t^2}{\sqrt{2}\pi\alpha}(1 - \frac{m_H^2}{m_\theta^2})^{1/2}. \tag{18}$$

This process has a rather large *QCD* correction.[13] If this is included, the right-hand side of Eq. (18) is multiplied by

$$\left(1 - \frac{4\alpha_s}{3\pi} a(m_H^2/m_\theta^2)\right)$$

Here, $a(x) \sim 10$ for $x \leq 0.8$, so that the correction reduces the naive rate. The branching ratio is reasonably large, but it is important to recall that the production rate for toponium in e^+e^- annihilation is not large. A 80 *GeV* toponium state has a production cross section of order 0.1 nb.

The product of a Higgs in hadron-hadron collisions occurs via several mechanisms. Since the Higgs coupling to light quarks is very small, the production of Higgs bosons from the annihilation of light quark-antiquark pairs strongly suppressed. There are too few heavy quarks inside the proton for their annihilation to generate a reasonable rate. There are two important mechanisms. Firstly, the Higgs can be produced via gluon-gluon fusion according to the Feynman diagram shown in Fig. 8.[14] This graph contains a vertex coupling the Higgs to a quark-antiquark pair, which is proportional to the quark mass. Consequently, the rate from this process depends sensitively upon the mass of the top quark. The production rate at center-of-mass energy $\sqrt{s}$ is a proton-proton collision is given by

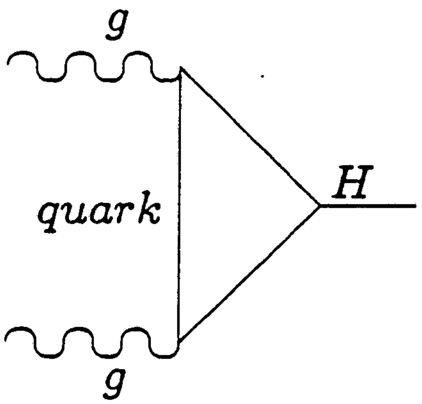


Figure 8: Feynman diagram showing the production of a Higgs boson via gluon-gluon fusion.

$$\sigma(pp \to H + X) = \frac{G_F \pi}{32\sqrt{2}} \left(\frac{\alpha_s}{\pi}\right)^2 \eta^2 \int_{M_H^2/s}^{1} \frac{M_H^2}{s} g(x) g(m_H^2/sx) dx. \qquad (19)$$

$g(x)$ is the gluon distribution of a proton. Defining $\epsilon_i = 4m_i^2/m_H^2$ for a quark of mass m_i, η is given by

$$\eta = \sum_i \frac{\epsilon_i}{2}(1 + (\epsilon_i - 1)\phi(\epsilon_i))$$

with

$$\phi(\epsilon) = \begin{cases} -[\sin^{-1}(1/\sqrt{\epsilon})]^2 & \epsilon > 1 \\ \frac{1}{4}(\log(\eta_+/\eta_-) + i\pi)^2 & \epsilon < 1 \end{cases}$$

where $\eta_\pm = 1 \pm \sqrt{1-\epsilon}$.

An alternative mechanism is shown in Fig. 9.[15] At large values of m_H, the rate from this mechanism becomes large due to the large width for $H \to WW$. The exact formula for this rate is complicated; it simplifies drastically in the so-called effective W approximation. This approximation assumes that the W's are emitted parallel to the incoming quarks and they are treated as if they are on mass-shell. It is similar to the effective photon approximation used to describe two-photon reactions in e^+e^- annihilation where the electron beams are treated as sources of

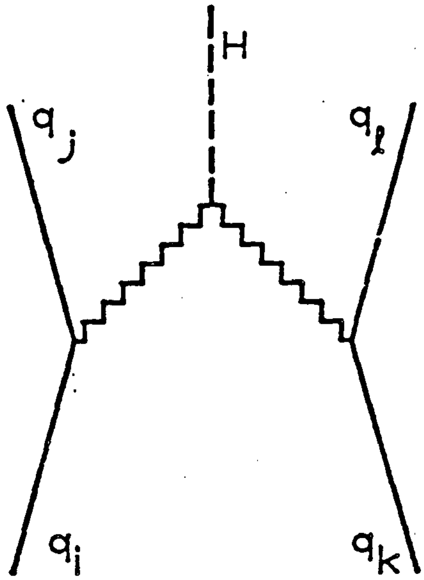


Figure 9: Feynman diagram showing the process $qq \rightarrow H + qq$.

on-shell photons. In this approximation the cross-section for $q + q \rightarrow H + qq$ via intermediate W's is given by[16]

$$\sigma\,(q_i + q_j \rightarrow H + q_j + q_i) = \frac{1}{16 M_W^2}\left(\frac{\alpha}{\sin^2\theta_W}\right)^3$$
$$\times\,[(1 + m_H^2/\hat{s})\log(\hat{s}/m_H^2) - 2 + 2m_H^2/\hat{s}]\theta(-e_i e_j) \tag{20}$$

where $\sqrt{\hat{s}}$ is the center-of-mass energy of the qq system and e_i is the charge of quark of type i. This may be converted into a hadronic cross-section via the parton model. In the case of intermediate Z bosons the factor $\theta(-e_i e_j)$ is replaced by $\frac{1}{\cos\theta_W^6}(v_i^2 + a_i^2)(v_j^2 + a_j^2)$ where v_i and a_i are given by,

$$a_i = T_3$$
$$V_i = T_3 - 4Q\sin^2\theta_W$$

where T_3 is the weak isospin of a fermion of charge Q.

This mechanism will only be important at the SSC; cross-sections evaluated at Tevatron and Sp$\bar{\text{p}}$S energies are dominated by the gluon fusion process. The rates for Higgs production are shown in Fig. 10. There are other mechanisms leading to final states with $H + Z, W + Z$ and $H + t\bar{t}$.[17,18]

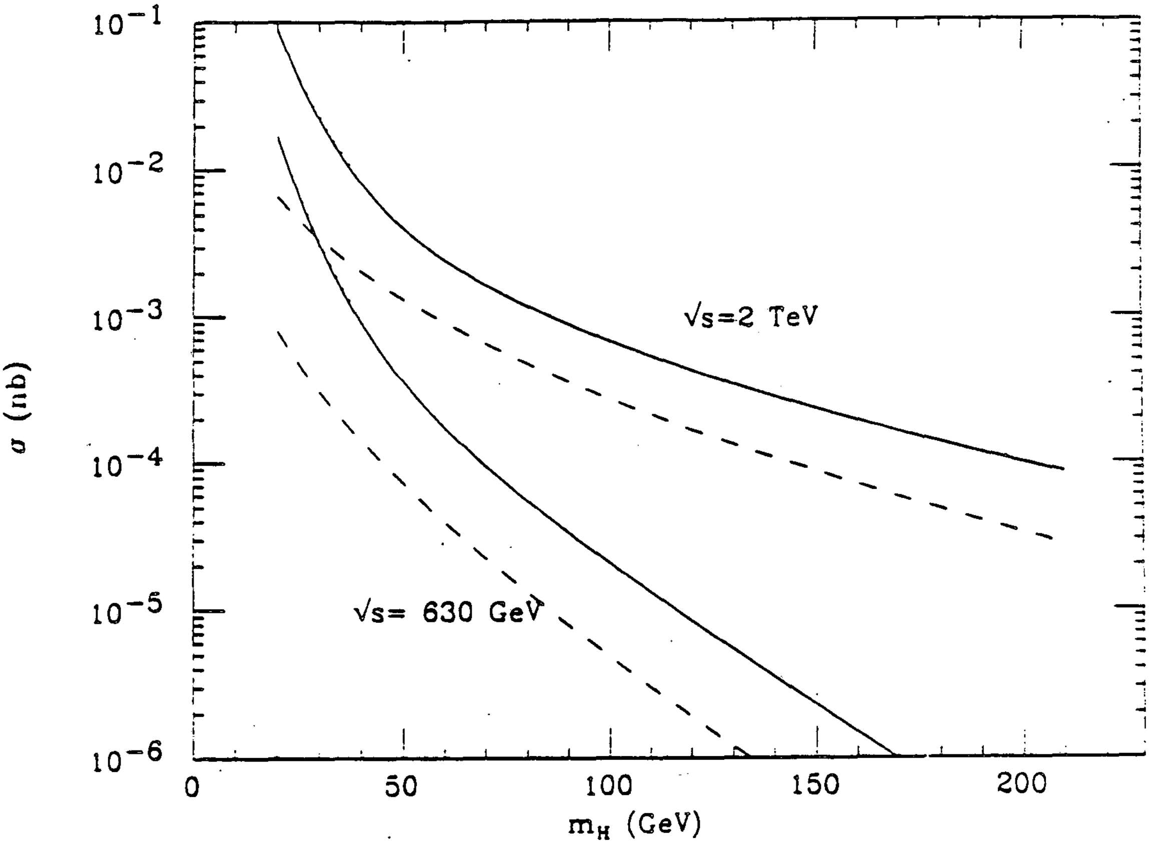


Figure 10: The cross-section $p\bar{p} \to H + X$ as a function of Higgs mass. The solid (dotted) lines correspond to the gluon fusion process of figure 25 with a top quark mass of 150 (40) GeV , and the dashed to the WW fusion process of figure 26.

A comparison of these rates at the SSC is shown in Fig. 11. At the Tevatron the rates are reasonable only for Higgs masses less than 150 GeV or so. In this mass region the Higgs will decay dominantly to $t\bar{t}$ if the t quark is light enough. There is a large background from the QCD production of $t\bar{t}$ pairs which will make detection difficult even if the t quark can be identified efficiently.

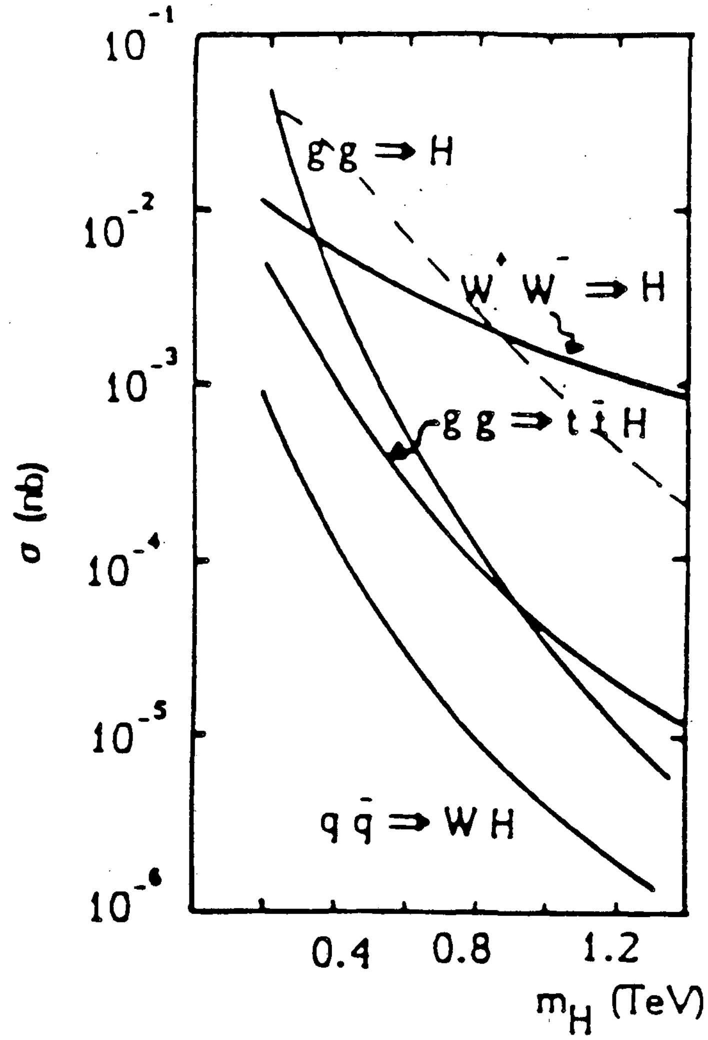


Figure 11: A comparison of the various Higgs production mechanisms in pp collisions at $\sqrt{s}$ =40 TeV. A top quark mass of 40 GeV is assumed for the solid lines and 150 GeV for the dashed.

If the $t\bar{t}$ channel is not open, the Higgs will decay to $\tau^+\tau^-$ with a branching ratio of $m_\tau^2/3m_b^2 \sim 4.5\%$. The only background source of τ pairs is Drell-Yan production $p\bar{p} \to \tau^+\tau^- + X$, via a virtual Z or photon. Figure 12 shows the signal and background in the τ pair channel. I have assumed a resolution of 10 *GeV* in the $\tau^+\tau^-$ invariant mass. It can be seen that the signal to background ratio is rather poor. This figure assumes a top quark mass of 150 GeV. The tau final state can be identified from the one-prong tau decays ($\tau \to e\nu\nu$, $\mu\nu\nu$, $\pi\nu$, etc.). Energy is lost into neutrinos so that the resolution in the $\tau^+\tau^-$ invariant mass will be poor. The experiment is clearly very difficult.

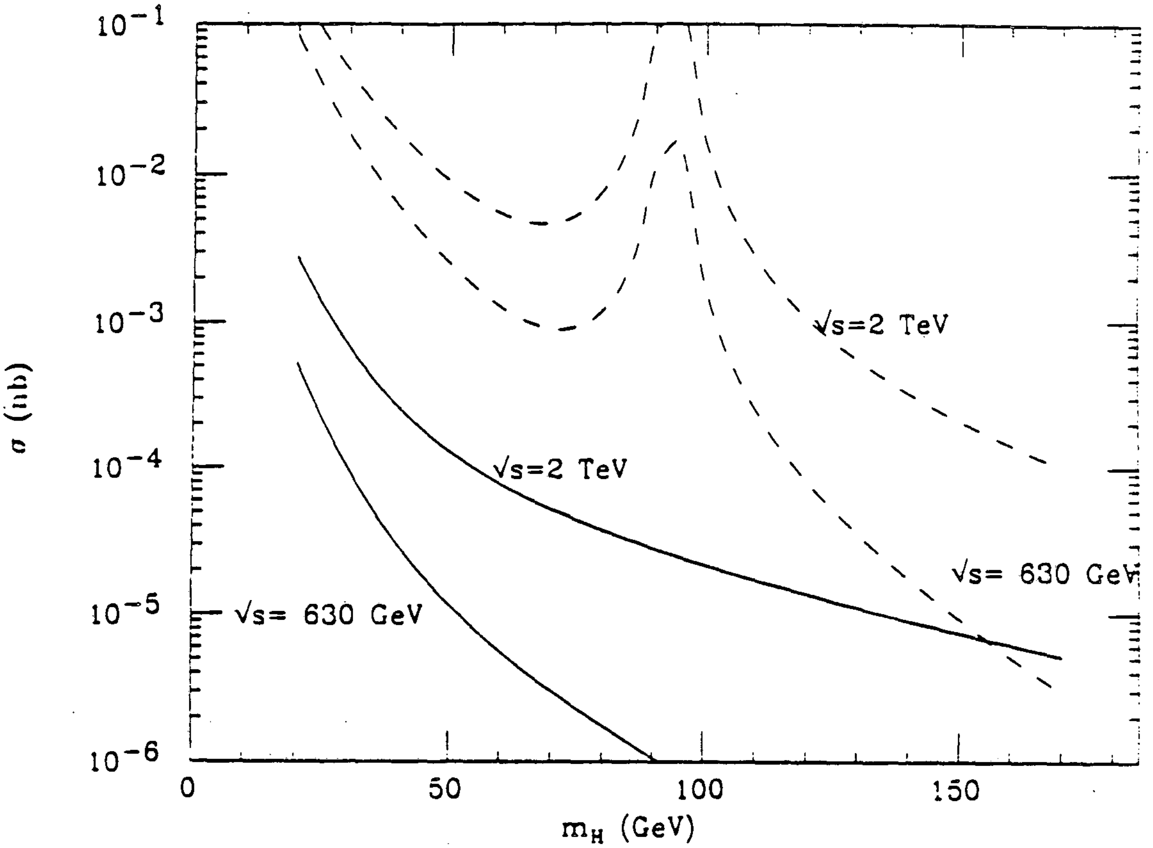


Figure 12: A comparison of the signal and background for the process $p\bar{p} \to H + X \to \tau^+\tau^- + X$. It is assumed that $m_H < 2m_t$. The background is calculated from the Drell-Yan process (see Sect. 3) being $d\sigma/dM \;\; \Delta M$. The resolution in the invariant mass of the tau pair (ΔM) is taken to be 10 GeV.

There is one other possibility. The Higgs can decay to two photons with the branching ratio[21]

$$BR(H \to \gamma\gamma) \sim \frac{m_H^2 \, |A| \, \alpha^2}{6\pi^2 m_b^2}. \tag{21}$$

I have assumed that $m_t > m_H/2$. Here A is a number arising from the W, quark and lepton loop diagram. Its value depends upon the masses of the particles involved but it is of order four for Higgs masses around 100 GeV. The background arises from the production of photon pairs via quark-antiquark collisions and is not too large. Unfortunately the branching ratio is so small that there are insufficient events for this decay mode to be useful. It has been suggested[22] as a possible mode at the SSC where the event rates are much larger.

What can we conclude about the prospects for finding the Higgs in the near future? If its mass is less than 40 GeV or so, it should be found in the decay of the Z either at the SLC or at LEP. Masses larger than this can be probed in the decay of toponium, if toponium exists in an accessible mass range. Notice that if $m_t > m_b + M_W$, the top quark will decay too quickly for narrow toponium bound states to exist. Higgs masses up to 100 GeV can be probed in the early 1990's at LEP when the energy is increased to 100 GeV per beam. Higgs bosons of mass greater than this will have to wait for the SSC.

It is convenient to distinguish two range of Higgs mass. If $m_H > 2M_W$, the Higgs boson will decay dominantly into W or Z pairs. In this case there is a background from the production of a W or Z pair via quark–antiquark annihilation. Figure 13 shows the rate for $pp \to H+X \to ZZ+X$ where the Z's are required to

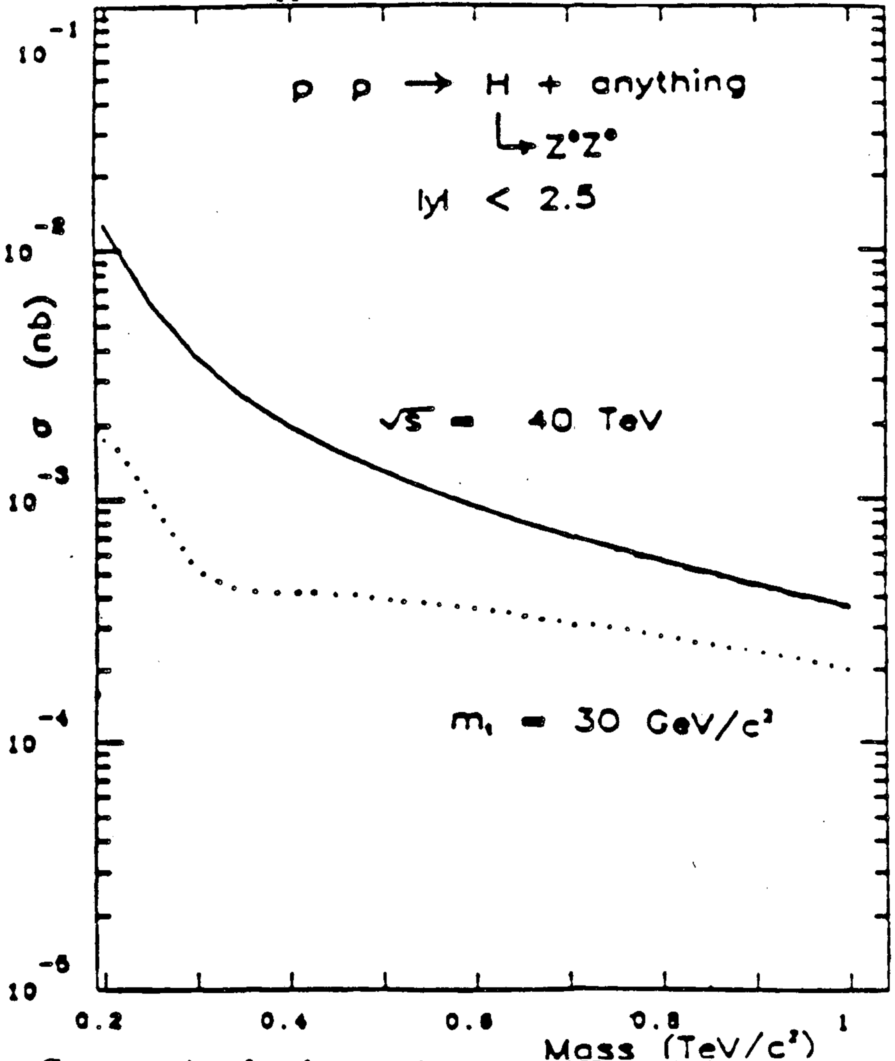


Figure 13: Cross–section for the reaction $pp \to (H \to ZZ) + X$ with $m_t = 30$ GeV. The Z's must have rapidity $|y| < 2.5$. Also shown (dotted line) is $\Gamma d\sigma(pp \to ZZ + X)/dM$ where $M = M_H$ is the invarient mass of the Z pair and $\Gamma = \max(\Gamma_H, 10 \text{ GeV})$.

have rapidity $|y| > 2.5$. The background is estimated from $\frac{d\sigma}{dM}\Delta M$, where $\frac{d\sigma}{dM}$ is the cross section for the production of a Z pair of invariant mass or the resolution Δm is taken to be the larger of 10 GeV or the Higgs width. It is clear that background is not a problem. However, the cross–section is not large particularly if the Z bosons cannot identified efficiently. If both Z's decay to e^+e^- or $\mu^+\mu^-$ we have a combined branching ratio of .36%, and an experiment is feasible only if it has an integrated luminosity of $10 fb^{-1}$, which corresponds to the design luminosity of the SSC. It is possible to look at the decay $H \rightarrow ZZ \rightarrow e^+e^- + \nu\nu$, where one of the Z's decays to $\nu\bar{\nu}$. A missing momentum measurement can be used to determine the transverse momentum of this Z. One can then look at the distribution in transverse mass of the Z pair. The signal and background are shown in Fig. 14. This method is complementary to that involving four charged leptons and is most useful at large value of m_H where the large Higgs width results in poor resolution in the charged lepton channel.

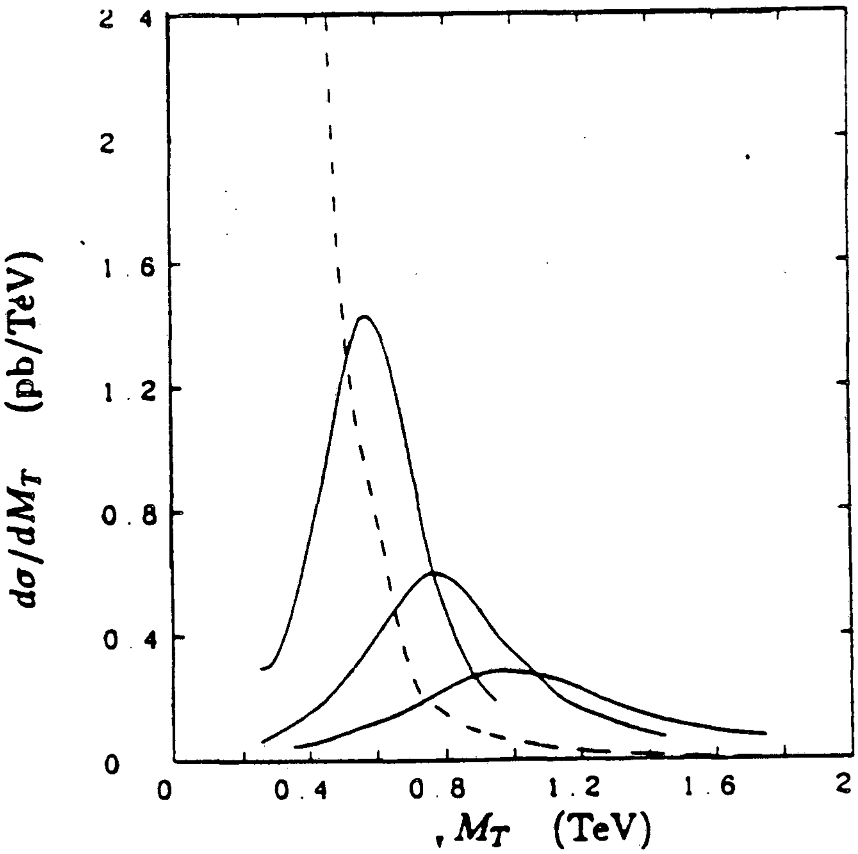


Figure 14: The tranverse mass distribution for $pp \rightarrow (H \rightarrow ZZ \rightarrow e^-e^+ + \nu\bar{\nu}) + X$ for various values of m_H at $\sqrt{s} = 40$ TeV in pp collisions. The

More events are available if one of the W's or Z's is observable in its hadronic decay modes. In this case there is a background from $W + Z$ jets. Attempts have been made to devise cuts which can reduce the background; there are some grounds for optimism.[20]

The situation is much more difficult when one considers the case $m_H < 2M_W$. If $m_H < 2m_t$, then the Higgs decays dominantly its top quark pairs. There is an enormous background from top quark pairs produced by gluon fusion; the signal is swamped even if top quarks can be identified perfectly. The final states $Z+H$ and $W+H$ offer some hope.[18] The background now arises from Z (or W)$+t\bar{t}$ and falls below the signal. However, it is still essential to identify the $t\bar{t}$ pair and reconstruct the Higgs. A detailed study revealed that this was almost impossible.[24] Firstly, the top quarks are not moving rapidly in the Higgs rest frame, so the jets from the $t\bar{t}$ pair overlap making it difficult to reconstruct the top quark momentum. Secondly, the probability of a semi–leptonic decay in the top jet is high so that energy is carried off by neutrinos and the resolution impaired. Thirdly, there is a background from $W + t\bar{b}$. We know of no way to find a Higgs at the SSC if its mass is less then $2M_W$ and greater than $2m_t$.

The situation is improved considerably if $m_H < 2m_t$.[26] Some of the rare processes discussed above may be used but the final state W (or Z) $+H$ now looks more promising. In the decay $H \rightarrow b\bar{b}$ all of the problems alluded to above in $H \rightarrow t\bar{t}$ are ameliorated. Figure 15 shows the signal and background for this channel.[26] Cuts have been applied requiring that there be a lepton with P_T relative to a quark jet of more than 1 GeV, and that jet be narrow. This reduces the background from light quark or gluon jets faking a b jet. Clearly the signal to background ratios would improve if a vertex detector were available to tag the b quark from its lifetime.

In conclusion, a Higgs boson of mass less than 80 GeV can probably be found at the LEP or the SLC. A heavier boson will have to wait for the SSC unless we are very lucky. What happens in variants of the standard model? The most popular variant is a supersymmetric model. In such a model there are two Higgs doublets, consequently there are three neutral and one charged Higgs particles. The coupling of some of the neutral particles to quarks is enchanced by a factor of v_1/v_2, the ratio of the vacuum expectation values of the two doublets. This coupling can enhance the production rate in hadron colliders via the gluon-gluon fusion process. However the coupling to W pairs does not increase with the Higgs mass in this case. Indeed there is a maximum value of about 10 GeV for the width to WW.[27] This implies that the WW fusion process will not be important for the production of such Higgs bosons. Fortunately, in these supersymmetric models, it is very likely that one of the neutral bosons is lighter than the Z,[28] so that it will be seen in e^+e^- annihilation.

> 1 GeV/c.

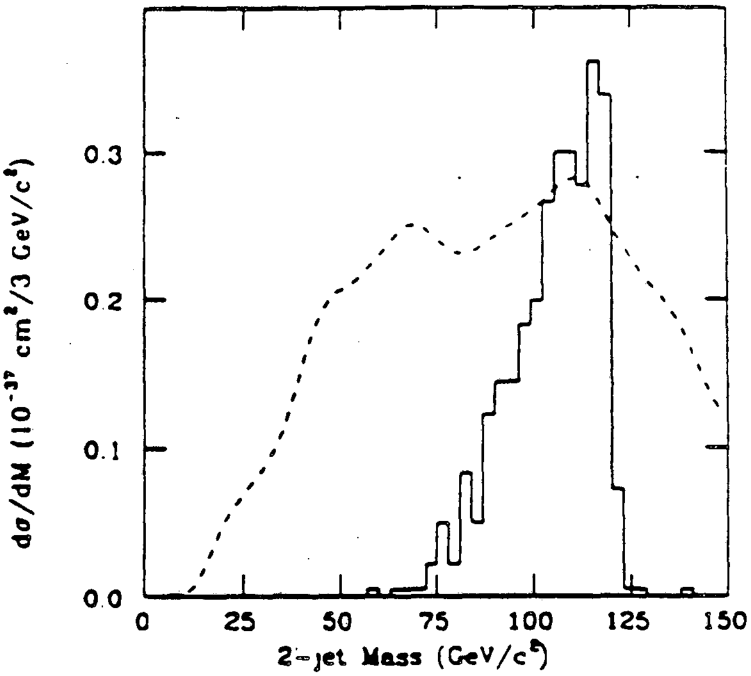


Figure 15: Two get mass spectrum arising from $pp \to W + (H \to \text{jet} + \text{jet}) + X$ and background for $pp \to W + 2$ jets as calculated in the PYTHIA Monte-Carlo. Experimental losses and resolutions we included. One jet is required to have a lepton with $p_T > 1$ GeV and both jets are required to have 70% if their energy in a cone of size $\Delta R = \sqrt{\Delta\phi^2 + \Delta\eta^2}$, where ϕ is the azimuthal angle and η is rapidity. The Higgs has mass 110 GeV.

ACKNOWLEDGMENTS

The work was supported by the Director, Office of Energy Research, Office of High Energy Physics, Division of High Energy Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098.

REFERENCES

1. S. Glashow, *Nucl. Phys.* 22:579 (1961);
   S. Weinberg, *Phys. Rev. Lett.* 19:1264 (1967);
   A. Salam, *in*: "Elementary Particle Theory," W. Svartholm, ed., Almquist and Wiksell, Stockholm (1968).
2. S. Coleman and E. Weinberg, *Phys. Rev. D* 7:1888 (1973).
3. S. Weinberg *Phys. Rev. Lett.* 36:294 (1976);
   A. Linde, *JETP Lett.* 23:64 (1976).

4. A. Guth and E. Weinberg, *Phys. Rev. Lett.* 65:1131 (1980).
5. R. Flores and M. Sher, *Ann. Phys.* 168:95 (1983).
6. C. Quigg, B.W. Lee and H. Thacker, *Phys. Rev. D* 16:1519 (1977).
7. M. Veltman, *Acta Phys. Polon. B* 8:475 (1977);
8. R. Peccei, these proceedings.
9. J.D. Bjorken, *in*: "Proceedings of the 1976 SLAC Summer Institute," M. Zipf, ed.
10. R.N. Cahn, M.S. Chanowitz and N. Fleishon, *Phys. Lett. B* 82:113 (1979).
11. B.L. Ioffe and V.A. Khoze, Leningrad report LINP-274 (1976).
12. F. Wilczek, *Phys. Rev. Lett.* 39:1304 (1977).
13. M. Vysotsky, *Phys. Lett. B* 97:159 (1980).
14. H.M. Georgi et al., *Phys. Rev. Lett.* 40:692 (1978).
15. R.N. Cahn and S. Dawson, *Phys. Lett. B* 136:196 (1986);
    S. Petcov and D.R.T. Jones, *Phys. Lett. B* 84:660 (1979).
16. M. Chanowitz and M.K. Gaillard, *Phys. Lett. B* 142:85 (1986).
17. Z. Kunszt, *Nucl. Phys. B* 247:339 (1984).
18. J.F. Gunion et al., *Phys. Rev. Lett.* 54:1226 (1985).
19. E. Eichten et al., *Rev. Mod. Phys.* 56:579 (1986).
20. J.F. Gunion and M. Soldate, *Phys. Rev. D* 34:826 (1986).
21. J. Ellis, M.K. Gaillard and D.V. Nanopoulos, *Nucl. Phys. B* 106:292 (1976).
22. G. Kane University of Michigan preprint UM TH 86-16 (1986), F. Gilman, J. Gunion and P. Wiesman, In preparation.
23. R. N. Cahn and M. Chanowitz, Phys. Rev. Lett. 56 1237 (1986).
24. F. Gilman and B. Cox, in Proc. of 1986 Snowmass Summer Study ed.R. Donaldson and J. Morfin (FNAL, 1985).
25. F. Gilman and L. Price SLAC-PUB-4138 (1986).
26. H. Haber and J. Gunion, Nucl. Phys, B278 449 (1986).
27. P. Roy Phys. Rev. D33 2674 (1986).