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PHENOMENOLOGY OF THE HIGGS BOSON

by

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ABSTRACT

The phenomenology of the standard Weinberg-Salam Higgs boson is reviewed with particular emphasis on production mechanisms in high energy e^+e^- and hadron-hadron collisions. The production processes relevant for the ISABELLE and TEVATRON energies are discussed and their backgrounds estimated. It is argued that the toponium production and radiative decay provides the most hopeful reaction to detect a Higgs in both the e^+e^- and the hadron-hadron machines.

I. INTRODUCTION

The problem of understanding the mechanism of mass generation is perhaps the most fundamental problem in elementary particle physics. Closely associated with it is the nature of the weak interaction scale, namely, why is the Fermi coupling constant $G_F = 1.05 \times 10^{-5}\, m_p^{-2}$?

In the standard theory of electroweak interactions[1] the masses of the fermions and the gauge bosons, which mediate the weak interactions, are governed by an order parameter, $\langle\phi\rangle_0$, the vacuum expectation value of an elementary scalar, color- and charge-neutral particle, ϕ. The mechanism which brings about $\langle\phi\rangle_0 \neq 0$ is now folklore and goes under the mystical name of spontaneous symmetry breaking. To be precise, one has a doublet of scalar Higgs fields $\phi = \binom{\phi^+}{\phi^0}$ having $SU(2)_L \times U(1)$ invariant couplings with the fermions and gauge bosons, and the Higgs potential has the form

$$V(\phi) = -\mu^2|\phi|^2 + \lambda|\phi|^4 \qquad \mu^2 > 0 \tag{1.1}$$

In order that $V(\phi)$ has a minimum for finite $|\phi|^2$, one must have $\lambda > 0$. Minimizing the potential (1.1) one finds that the minimum is not at $\phi = 0$ but at a non-zero point determined by μ and λ.

$$|\phi|^2 = \mu^2/\lambda \equiv v^2/2 \tag{1.2}$$

The situation that $v \neq 0$ is what is meant by having a spontaneously broken symmetry. Now if one rescales the field $\phi \to \phi - v$, then the resulting theory can be arranged in such a way that the fields ϕ^+, $\phi^- = \bar{\phi}^+$ and $\chi = 1/\sqrt{2}\, i\, (\phi^0 - \bar{\phi}^0)$ disappear from the lagrangian and these degrees of freedom become the longitudinal components of W^+, W^- and Z^0, respectively, which in turn become massive. This is what goes under the name of Higgs mechanism. The surviving scalar field, which we shall call H, has the form

$$H^0 \equiv \frac{\phi^0 + \bar{\phi}^0}{\sqrt{2}} - v \qquad (1.3)$$

H is a physical field. Except mass, it has all the attributes of the vacuum. The reshuffled lagrangian has the form

$$\mathcal{L}(W,f) = V(H) + g_{fH}\, \bar{f} f H^0 - \frac{1}{4} g^2 W_L^+ W_L^- (H^0)^2 - g_{WH}(W_L^+ W_L^- H^0) + g_{ZH}\, Z_L^2\, H^0 + \ldots\ldots \qquad (1.4)$$

where the parameters g_{fH}, g_{WH} are completely determined in terms of the masses m_f, m_W, m_Z and the vacuum expectation value v. One has

$$g_{fH} = m_f/v = (\sqrt{2}\, G_F)^{1/2}\, m_f$$
$$g_{WH} = 2m_W^2/v = 2(\sqrt{2}\, G_F)^{1/2}\, m_W^2 \qquad (1.5)$$

The proportionality of the Higgs-fermion-fermion coupling and Higgs-Boson-Boson couplings to their masses is no accident. It follows from demanding that the Higgs couplings induce fermion and boson masses after spontaneous symmetry breaking.

The difficulty in having a clean signal in the search for H^0 is precisely quantified through Eqs. (1.4), namely that the Higgs likes to couple to the fermions and bosons with an intrinsic strength proportional to their masses. Thus, the production rate in all lepton (e, μ, ν_e, ν_μ)-induced reactions is either zero or miniscule. In the same spirit the <u>tree level</u> couplings involving photons $g_{\gamma\gamma H}$ and the gluon g_{GGH} are zero, which is a consequence of gauge invariance. Thus $g_{\gamma\gamma H}$ (and g_{GGH}) are non-zero only in higher orders in α (and α_s), consequently the γ-induced and gluon-induced reactions are

also intrinsically very small. On the other hand the decays of the Higgs are dominantly into heavy fermions and bosons, leading to multiparticle final states thereby making the traditional $\mu^+\mu^-$ invariant mass searches prohibitively small.

The purpose of this talk is to review attempts in producing and detecting the Higgs boson using high energy machines already available and being planned. While the phenomenology of the Higgs boson will be reviewed in general, I will concentrate more on production mechanisms in the hadronic reactions which are more relevant for the purpose of the present meeting. The other reason is that the opportunities that high energy e^+e^- machines, like LEP, provide in Higgs searches have already been emphasized in the literature.[2] The role of hadron machines in Higgs searches has not received the attention it deserves. I will try to convince you that the intrinsic production cross sections for the Higgs in high energy proton-proton and proton-antiproton colliding machines are not small but the detection of the signal needs new thinking and strategies on the part of our experimental colleagues.

The plan of this talk is as follows. I will start by reviewing the bounds on and an estimate of the Higgs mass. I shall then discuss the various decay mechanisms. The production mechanisms of the Higgs boson is the main subject of this talk and I shall review them emphasizing in particular the radiative decay of the toponium,[3] $\theta_t \to H^0 + \gamma$ and how to make use of this mechanism in hadronic processes. The radiative decay of toponium in my opinion provides the best chance of observing a Higgs both in the e^+e^- machines and at ISABELLE and TEVATRON, if the masses of the toponium and Higgs permit such a decay. Wherever relevant I shall compare the Higgs scenario with the hypercolor scenario of dynamical symmetry breaking,[4] which also admits (almost) point-like light (pseudo-)scalar particles, though this subject will be reviewed by Baqi Beg[5] and Gordy Kane[6] in separate talks in this meeting.

II. MASS OF THE WEINBERG-SALAM HIGGS BOSON UPPER BOUND ON m_{H^o}

The mass of the Higgs boson is in general not determined by the theory since it depends on the unknown quartic coupling constant λ in the Higgs potential (1.1). However an upper bound on m_{H^o} can be obtained in terms of an upper bound on λ. The coupling constant λ is bounded by $\lambda < 1$, otherwise the perturbation theory in λ breaks down. The precise bound is obtained if one considers the scattering of longitudinally polarized $W^\pm$ bosons: $W_L^+W_L^- \to W_L^+W_L^-$.[7] The amplitude via the γ and Z exchange is linearly divergent!

$$T^{\gamma,Z} \underset{s\to\infty}{\to} \frac{G_F}{\sqrt{2}} s(1+\cos\theta);\; s = 4E_{W^\pm}^2 \tag{2.1}$$

This linear divergence (in s) is cancelled by the Higgs contribution giving

$$T^{\gamma,Z,H^o} \sim -\left(\frac{4G_F}{\sqrt{2}}\right) m_H^2 \tag{2.2}$$

Using the partial wave decomposition, one has

$$T = 16\pi \sum_J (2J+1)\, t^J P_J(\cos\theta) \tag{2.3}$$

Unitarity bound for each partial wave is

$$|t^J| \leq 1 \tag{2.4}$$

For $J = 0$, this implies $T \leq 16\pi$ which translates to an upper bound on m_H^2.

$$m_H^2 \leq 4\pi\sqrt{2}/G_F \tag{2.5}$$

The unitarity bound (2.5) could be refined by considering the 3-channel coupled system consisting of $W_L^+W_L^-$, $1/\sqrt{2}\, Z_LZ_L$ and $1/\sqrt{2}\, H^oH^o$ which gives the bound

$$m_H^2 \leq \frac{8\pi\sqrt{2}}{3G_F} = 1\ \text{TeV}^2 \tag{2.6}$$

which is disappointingly large! If it turns out that nature has chosen a value for m_{H^o} close to its upper bound, then the Higgs boson is beyond the reach of all present and planned machines like LEP, ISABELLE and TEVATRON!

LOWER BOUND ON m_{H^o}

The lower bound on m_{H^o} comes by considering the radiative corrections to the Higgs potential (1.1). The 1-loop $[SU(2)_L \times U(1)]$ radiative correction gives the result

$$V(\phi) = -\mu^2|\phi|^2 + |\phi|^4 \ln(|\phi|^2/M^2) \times \left(\sum_{V=Z,W^\pm} 3m_V^4 + m_H^4 - 4\sum_f m_f^2\right) \tag{2.7}$$

where M is a mass parameter to absorb all $|\phi|^4$ terms in $V(\phi)$. The parameters μ and M are so chosen so that the vacuum is stable, i.e. $V(v^2/2) < V(0)$ where v is determined by demanding

$$\left.\frac{\partial V\text{rad.}}{\partial\phi}\right|_{|\phi| = v/\sqrt{2}} = 0 \tag{2.8}$$

In the presence of radiative corrections, there is the amusing possibility that one could have $\mu^2 < 0$, and still arrange spontaneous symmetry breaking. However, in that case there will be several minima. It is conceivable that the theory (and the universe) is at a local minimum and will decay to the absolute minimum leading to catastrophic consequences. The rate of such a transition depends critically on m_{H^o}. It has been shown by Linde[8] that if $m_{H^o} > 260$ MeV, then the rate of this transition would be so slow that it is not of any immediate worry!

$$t(\phi_{min} = v/\sqrt{2} \to \phi_{min} = 0) > 10^{10}\text{yrs for } m_{H^o} > 260\ \text{MeV} \tag{2.9}$$

Of course, if m_{H^o} is close to the Linde bound, it will have important cosmological consequences.

Demanding $\varphi_{min} \neq 0$ gives an upper bound on $(-\mu^2)$ which translates to a lower bound on m_{H^o}.[9]

$$m_H^2 = \left.\frac{\partial^2 V}{\partial\phi^2}\right|_{\phi = v/\sqrt{2}} > \frac{3}{16\pi^2 v^2} \sum_{V=W^\pm,Z} m_V^4 > (7.2\ \text{GeV})^2 \quad \text{for } \sin^2\theta_W = 0.20 \tag{2.10}$$

Thus, the present bounds on m_{H^o} are

$$1\ \text{TeV} > m_{H^o} \geq 7.2\ \text{GeV} \tag{2.11}$$

COLEMAN-WEINBERG ESTIMATE OF m_{H^o}

E. Weinberg and S. Coleman[10] have pointed out that one could set $\mu^2 = 0$, $\lambda > 0$ and achieve spontaneous symmetry breaking via radiative corrections. It will then fix the mass of the Higgs boson to be

$$m_H^2 = \frac{3\alpha^2 v^2}{8} \left\{ \frac{(2 + \sec^4\theta_W)}{\sin^4\theta_W} - O(m_f/m_H)^4 + O(\alpha) \right\} \quad (2.12)$$

which gives (neglecting fermion mass and $O(\alpha)^3$ contributions) (v = 247 GeV)

$$m_{H^0} \approx 10.4 \text{ GeV}^{+0.5 \text{ GeV}}_{-0.4 \text{ GeV}} \text{ for } \sin^2\theta_W = 0.2^{+0.01}_{-0.01} \quad (2.13)$$

Thus, for the present value of the Weinberg-Salam-Glashow angle $\sin^2\theta_W = 0.215$ (2.12) predicts $m_{H^0} \gtrsim 11$ GeV, which puts it above the mass of the heaviest observed bound ($Q\bar{Q}$) system Υ, Υ', Υ'' .... It should perhaps be remarked that the heavy fermion mass contribution decreases the estimate (2.13) for m_{H^0} (though $\Delta m_H(f) < 10$ MeV for $m_f \leq 20$ GeV) and that the estimate of m_{H^0} is based on one-loop calculations. The possibility (2.13) is very exciting from the point of view of the proposed e^+e^- and proton machines and I shall review the consequences of a Higgs with $m_H \gtrsim 11$ GeV in this talk.

III. DECAYS OF THE WEINBERG-SALAM HIGGS BOSON

The decays of the Higgs boson are determined by the couplings in (1.4) and (1.5). There are no tree level couplings $g_{H^0\gamma\gamma}$ or g_{H^0GG}. However, these couplings are induced at the 1-loop (QFD) level. Thus, the decays of H^0 up to 1-loop level are:

$$\begin{aligned} H^0 &\to Z^0Z^0 \\ &\to W^+W^- \\ &\to \ell^+\ell^- \qquad \ell = e, \mu, \tau \\ &\to q\bar{q} \qquad q = u, d, s, c, b, t \\ &\to GG \\ &\to \gamma\gamma \end{aligned} \quad (3.1)$$

where the decays into weak gauge bosons are allowed only if $m_{H^0} > 2m_{W^\pm}$, $2m_{Z^0}$, in which case the H^0 decays would be totally dominated by these modes. The decay widths into Z^0Z^0 and W^+W^- are given by

$$\Gamma(H^0 \to W^+W^-) = \frac{G_F M_W^2}{8\pi\sqrt{2}}\, m_H \frac{(1 - x)^{1/2}}{x} (3x^2 - 4x + 4) \quad (3.2a)$$

$$\Gamma(H^0 \to Z^0Z^0) = \frac{G_F M_W^2}{8\pi\sqrt{2}}\, m_H \frac{(1 - x')^{1/2}}{x'} (3x'^2 - 4x' + 4) \quad (3.2b)$$

where

$$x = \frac{4M_W^2}{m_H^2}, \quad x' = \frac{4M_Z^2}{m_H^2} = \frac{x}{\cos^2\theta_W}$$

For $m_{H^0} \geq 200$ GeV, $\Gamma_{H^0} > 1$ GeV and the width increases very fast becoming bigger than m_{H^0} for $m_{H^0} \geq 1$ TeV. In that limit weak interactions become strong and the perturbation theory would break down, as I have already remarked. Leaving the question of producing a 200 GeV object apart, detecting such an object would be relatively easy through the modes

$p p(\bar{p}) \to H^0 + x$
  $\hookrightarrow Z^0Z^0$
    $\hookrightarrow \ell^+\ell^-$ ($\ell = e, \mu, \tau$)
    $\hookrightarrow$ hadrons, $\ell^+\ell^-$

and (3.3)

$p p(\bar{p}) \to H^0 + x$
  $\hookrightarrow W^+W^-$
    $\hookrightarrow \ell^-\bar{\nu}_\ell$
    $\hookrightarrow \ell^+\nu_\ell$, hadrons

Note that such an object will not be confused with the technieta,[11] the massive colored pseudoscalar object present in technicolor/hypercolor theories which would decay predominantly into a $t\bar{t}$ pair, having very different event topologies. The color neutral pseudo-Goldstone boson, PGB, π'^0, even if it is massive ($m_{\pi'^0} > 2m_{W^\pm}$) will not decay at the tree level through the modes

$$\pi'^0 \to Z^0Z^0,\ W^+W^- \quad (3.4)$$

The couplings $g_{\pi'^0 Z^0 Z^0}$ and $g_{\pi'^0 W^+ W^-}$ are nonzero only at the one-loop level and hence very small.[12] The scenario of a heavy Higgs boson, $m_{H^0} > 2m_{Z^0}$, while somewhat discomforting from the point of view of production will have the redeeming feature that it's detection will be easy and it would not be confused with heavier pseudo-scalar objects of the hypercolor scenario.

Let us now concentrate on the scenario in which $m_{H^0} < 2m_{W^\pm}$, in that case the decays of H^0 would be dominated by the heaviest fermion pair allowed by phase space. The decay widths for $H^o \to \ell^+\ell^-$ and $H^o \to q\bar{q}$ are given by

$$\Gamma(H^o \to \ell^+\ell^-) = \frac{G_F m_\ell^2}{4\sqrt{2}\,\pi}\, m_H \left(1 - \frac{4m_\ell^2}{m_H^2}\right)^{3/2}$$
$$\Gamma(H^o \to q\bar{q}) = 3\,\frac{G_F m_q^2}{4\sqrt{2}\,\pi}\, m_H \left(1 - \frac{4m_q^2}{m_H^2}\right)^{3/2} \qquad (3.5)$$

where the factor 3 for $H^o \to q\bar{q}$ is due to color.

Next we shall calculate the decays $H^o \to \gamma\gamma$,[13] and $H^o \to GG$.[3] These decays, which are allowed at the one-loop level, are quite amusing and if measured could be used as heavy quark counters. The effective $H^o \to 2\gamma$ coupling can be expressed as (see Fig. 1)

$$F(k_1, k_2, q) = \frac{2\alpha}{\pi v} I_{\mu\nu}(k_1,k_2)\, \varepsilon_\mu \varepsilon_\nu \qquad (3.6)$$

where gauge invariance dictates that $I_{\mu\nu}$ be expressed as

$$I_{\mu\nu} = (k_{1\mu}k_{2\nu} - k_1 \cdot k_2\, g_{\mu\nu})I$$

and we decompose I according to the contributions from the fermion- and gauge boson-loop

$$I = I_W + \sum_f I_f \qquad (3.7)$$

With this notation the decay width is given by

$$\Gamma(H^o \to 2\gamma) = \frac{G_F \alpha^2}{8\sqrt{2}\,\pi}\, m_H^3\, |I|^2 \qquad (3.8)$$

The two contributions to I can be calculated in a straight-forward way

$$I_W = \frac{-7}{4} + O(m_H^2/m_W^2) \qquad (3.9)$$

$$I_f = N_c Q_f^2 L_f \qquad (3.10)$$

where N_c is the color factor (=3), Q_f is the charge of the fermion in the loop and the factor L_f is given by

$$L_f = m_f^2 \int_0^1 dx \int_0^{1-x} dy\, \frac{1 - 4xy}{m_f^2 - m_H^2 xy} \qquad (3.11)$$

$$L_f \to \tfrac{1}{3}[1 + O(m_H^2/m_f^2)],\quad m_H^2/m_f^2 << 1$$

Thus, for large m_f, L_f becomes independent of m_f and counts the heavy fermion degrees. In particular, for a complete 3-generation family (e, μ, τ; u, d, s, c, b, t) we have

$$I(3) \equiv \sum_{3\ \text{gen.}} I_f = 0.7 \qquad (3.12)$$

if (N-3) generations of heavy fermions exist then

$$I(N) = I(3) + \tfrac{8}{9}(N-3) \qquad (3.13)$$

Putting (3.9) and (3.13) together, we get

$$\begin{aligned} I &\sim -1 && \text{for } N = 3 \\ &\sim -0.1 && \text{for } N = 4 \\ &\geq 0.8 && \text{for } N \geq 5 \end{aligned} \qquad (3.14)$$

Thus, in principle, $BR(H^o \to 2\gamma)$ is a very sensitive way to 'feel' a fourth generation of fermions. In practice, however

$$\begin{aligned} BR(H^o \to 2\gamma) &\sim 4 \times 10^{-5}\ |I|^2 \text{ for } 2m_D < m_{H^o} < 2m_B \\ &\sim 4 \times 10^{-6}\ |I|^2 \text{ for } 2m_B < m_{H^o} < 2m_T \end{aligned} \qquad (3.15)$$

which is hopelessly small.

Finally, we quote the result for the decay rate $\Gamma(H^o \to GG)$, for which only the heavy quarks contribute in the loop[3] (Fig. 2).

$$\Gamma(H^o \to 2G) = \frac{G_F N_F^2}{36\sqrt{2}\,\pi}\left(\frac{\alpha_s(m_H^2)}{\pi}\right)^2 m_H^3 \qquad (3.16)$$

where N_F = number of heavy flavors. The decay rate $\Gamma(H^o \to 2G)$ is larger than the rate $\Gamma(H^o \to \mu^+\mu^-)$ for $m_{H^o} \geq 3$ GeV. However, it is still much smaller than the decay rate into the heaviest fermion pair. In addition $H^o \to 2G$ mode lacks a reliable trigger both in e^+e^- annihilation and hadron-hadron collisions and we shall not discuss it any more.

Traditionally, the searches for new particles (J/ψ, Υ) have been very successful in the hadron-hadron collisions through the mode $pp \to V + x$, $V \to \mu^+\mu^-$.

However, because of the peculiar $Hf\bar{f}$ couplings the branching ratio for $H^o \to \mu^+\mu^-$ is miniscule. To orient ourselves we note that

$$BR(H^0 \to \mu^+\mu^-) \simeq 2 \times 10^{-3} \text{ for } 2m_D < m_{H^0} < 2m_B$$
$$\simeq 2 \times 10^{-4} \text{ for } 2m_B < m_{H^0} < 2m_\tau$$
$$\leq 10^{-5} \quad \text{ for } 2m_\tau < m_{H^0} < 2m_W \tag{3.17}$$

Thus, though the production of H^0 in high energy hadron-hadron machines may be large, as we shall see in the next section, searching a peak in the $\mu^+\mu^-$ invariant mass would be a disappointing enterprise. Hadron machines require a useful trigger other than $\mu^+\mu^-$.

IV. PRODUCTION OF THE HIGGS BOSON

The Higgs production processes that I would like to discuss in this talk are:

i. Decays of vector bosons involving a Higgs

ii. Higgs production in pp and $p\bar{p}$ collisions

iii. Higgs production in lepton-hadron processes

and iv. Higgs production in e^+e^- annihilation.

Let me start with the production mechanism (i).

i. DECAYS OF VECTOR BOSONS INVOLVING A HIGGS

The vector particles in whose decays a Higgs boson may be produced are the Z^0 boson and the bound heavy flavor vector meson, Υ and the yet to be discovered state toponium, J_T. The Higgs boson may also be produced in the decays of the charged vector bosons, $W^\pm$, however, I don't know how to produce $W^\pm$ copiously to detect branching ratios at the level of 10^{-4} which one expects in the decays $W^\pm \to H^0 W^\pm_{vir} \to \ell^\pm\nu_\ell,\ q\bar{q}$.[14]

HIGGS IN THE DECAYS OF Z^0

There are two decay modes which have been advocated in the literature. The radiative decay[15]

$$Z \to H^0 + \gamma \tag{4.1}$$

and the decay[16]

$$Z \to H^0 + Z^0_{vir} \to \ell^+\ell^- \tag{4.2}$$

In the standard model, the coupling $ZH^0\gamma$ is not allowed at the tree level.

At the one-loop level it can be calculated using the diagrams shown in Fig. 3. The matrix element for the process (4.1) can be written as

$$m = \varepsilon^\mu_{(\gamma)}\, m_{\mu\nu}\, \varepsilon^\nu_{(Z^0)} \tag{4.3}$$

Electromagnetic gauge invariance ($m_{\mu\nu}k^\nu = 0$) gives

$$m_{\mu\nu} = (k_\nu p_\mu - k\cdot p\, g_{\mu\nu})\, a \tag{4.4}$$

in terms of which the width can be expressed as

$$\Gamma(Z \to H^0 + \gamma) = \frac{E_\gamma^{\,3} a^2}{12\pi} \tag{4.5}$$

leading to

$$\Gamma(Z \to H^0\gamma)/\Gamma(Z \to \mu^+\mu^-) \simeq \frac{\alpha^2}{\pi^2 \sin^2\theta_W} \left(\frac{E_\gamma}{M_Z}\right)^3 A^2 \simeq 2.6 \times 10^{-5} \left(\frac{E_\gamma}{M_Z}\right)^3 A^2 \tag{4.6}$$

The constant A from Fig. 3 is estimated to be

$$A_{(\text{Fermion loop})} = \frac{1}{3\cos\theta_W}\left(1 - \frac{8}{3}\sin^2\theta_W\right)$$
$$\simeq 0.7$$
$$A_{(W\text{-boson loop})} \simeq -(4.9 + 0.3(m_H^2/m_W^2)) \tag{4.7}$$

leading to

$$\Gamma(Z \to H^0\gamma)/\Gamma(Z \to \mu^+\mu^-) \simeq 7.8 \times 10^{-5}\left(1 - \frac{m_H^2}{m_Z^2}\right)^3$$
$$\simeq 10^{-6} \times \left(1 + .17\,\frac{m_H^2}{m_Z^2}\right) \tag{4.8}$$

Even with this small rate one would have to look at

$$Z^0 \to H^0 + \gamma, \quad H^0 \to \ell^+\ell^- \tag{4.9}$$

Barbellini *et al.*[2] have calculated the background to the process (4.9) due to the decay mode

$$Z^0 \to \ell^+\ell^-\gamma \tag{4.10}$$

and it looks formidable. Even after demanding that the photon be recoiling against $\ell^+\ell^-$, the signal from (4.9) would be buried in the background. In e^+e^- annihilation, where one could demand a monoenergetic photon, there may be some outside chance of observing (4.9) but in pp and $p\bar{p}$ collisions it is simply undoable.

Next, we discuss the process (4.2). In principle one could also look at the mode

$$Z^o \to H^o + Z^c_{vir} \quad \llcorner\!\to q + \bar{q}$$

but its separation from the dominant Z^o decays $(Z^o \to q\bar{q})$ is not easy. The rate for (4.2) is given by[16]

$$\frac{1}{\Gamma(Z^o \to \mu^+\mu^-)} \frac{d\Gamma}{dX_H}(Z^o \to H^o + \ell^+\ell^-) = \frac{\alpha}{4\pi \sin^2\theta_W \cos^2\theta_W} \times \frac{1}{(X_H - m_H^2/m_Z^2)^2}\left[1 - X_H + \frac{X_H^2}{12} + \frac{2}{3}\frac{m_H^2}{m_Z^2}\right]\left[X_H^2 - \frac{4m_H^2}{m_Z^2}\right]^{1/2} \tag{4.11}$$

where $X_H = 2E_H/m_Z$. The rates for $Z^o \to H^o\ell^+\ell^-$ and $Z^o \to H^o + \gamma$ are shown in Fig. 4. For $m_{H^o} \simeq 10$ GeV it leads to a branching ratio

$$BR(Z^o \to H^o + \mu^+\mu^-) \simeq O(10^{-5}) \tag{4.12}$$

which certainly is observable in a high luminosity Z^o-factory. A very good handle on the process (4.2) is obtained if one looks at the shape of the dimuon distribution which peaks at large values of $m_{\mu\mu}$.[16]

$$\left.\frac{d\sigma}{dm^2_{\mu\mu}}(Z^o \to H^o\mu^+\mu^-)\right|_{\chi} = 0 \tag{4.13}$$

where $\chi = 0.95\, m_H/m_Z$. The background to (4.2) comes from the dominant decay mode

$$Z^o \to Q\bar{Q} \to \mu^+\mu^- + X \tag{4.14}$$

However, this peaks at the low invariant mass of the dimuons, thus providing a clear signal. The invariant mass distribution, shown in Fig. 5, is a characteristic of the Higgs mode (4.2) and is very different if one considers instead the process involving a pseudo-Goldstone boson.

$$Z \to \pi'^o + Z^c_{vir} \quad \llcorner\!\to \mu^+\mu^- \tag{4.15}$$

Of course the rate expected for $Z \to \pi'^o + \mu^+\mu^-$ is prohibitively small.[12]

$$\frac{\Gamma(Z^o \to \pi'^o\mu^+\mu^-)}{\Gamma(Z^o \to H^o\mu^+\mu^-)} < 10^{-4} \tag{4.16}$$

The decay (4.2) in a Z^o factory is one of the most promising places to observe a Higgs if $m_{H^o} \lesssim 40$ GeV, beyond which both the branching ratio becomes very small and the signal to background separation no longer remains that good. In pp and $p\bar{p}$ machines (4.2) is buried under the background.

$$pp(p\bar{p}) \to Z^o + X \quad \llcorner\!\to \mu^+\mu^- \tag{4.17}$$

which is at least three orders of magnitude bigger than for the process[17]

$$pp(p\bar{p}) \to Z^o + X \quad \llcorner\!\to H^o + \mu^+\mu^- \tag{4.18}$$

The situation becomes better if one looks for the H^o signal in trilepton final state via

$$pp(p\bar{p}) \to Z^o + X \quad \llcorner\!\to H^o + \mu^+\mu^- \quad \llcorner\!\to \mu^\mp + X' \tag{4.19}$$

which has been calculated in Ref. 18 and shown in Fig. 6. The cross section for (4.19) is of order 10^{-39} cm^2 at ISABELLE and I guess the background from the Drell-Yan process to (4.19) may swamp the feeble signal. The use of (4.19) for Higgs search is going to be a formidable task.

HIGGS IN THE DECAYS OF QUARKONIA

Since the couplings of the Higgs boson to a fermion pair favors heavy fermions [see (1.5)], it immediately suggests that the search of Higgs will be profitable in processes in which heavy fermions are involved. Wilczek[3] suggested that if $m_{H^o} < m_\Upsilon$ then it could be produced in the radiative decay

$$\Upsilon(9.46) \to H^o + \gamma \tag{4.20}$$

The decay rate for (4.20) can be calculated using non-relativistic quark model and one obtains[3]

$$\frac{\Gamma(T \to H^0 + \gamma)}{\Gamma(T \to \mu^+\mu^-)} = \frac{G_F m_T^2}{4\sqrt{2}\pi\alpha}\left(1 - \frac{m_H^2}{m_T^2}\right) \tag{4.21}$$

using the known branching ratio for $T \to \mu^+\mu^-$. (4.21) already gives an upper limit for the branching ratio $BR(T \to H^0 + \gamma) < 2.5 \times 10^{-4}$. It has been pointed out in Ref. 19 that if m_{H^0} is close to m_T, then the estimate (4.21) should be corrected to take into account the dipole nature of the radiative decay. Correcting for this we multiply the right hand side of Eq. (4.21) by the dipole factor K, where a phenomenological estimate for K is[19]

$$K = \frac{\frac{m_T}{2}\left(1 - m_H^2/m_T^2\right)}{\left[\frac{m_T}{2}\left(1 - m_H^2/m_T^2\right) + \Lambda\right]} \underset{m_H^2/m_T^2 \ll 1}{\longrightarrow} 1 \tag{4.22}$$

and Λ is an onium potential dependent factor $\sim O(1\ \text{GeV})$. The K factor will further reduce the branching ratio depending upon the mass difference $(m_T - m_{H^0})$. The possibility (4.20) will soon be checked at CESR and DORIS. However, if the Coleman-Weinberg estimate of $m_{H^0} \gtrsim 11$ GeV is correct, then the decay (4.20) is not accessible. However the decay $J_T \to H^0 + \gamma$ would have a large branching ratio in the decay of the (yet to be discovered) toponium state, J_T. The present experimental limit on the mass m_{J_T} from PETRA experiments is[20]

$$m(J_T) > 36.7\ \text{GeV} \tag{4.23}$$

implying

$$\frac{\Gamma(J_T \to H^0 + \gamma)}{\Gamma(J_T \to \mu^+\mu^-)} > 0.13 \tag{4.24}$$

which is not a small number. On the other hand if the top quark mass is larger than $m_Z(\simeq 93$ GeV), then J_T would decay dominantly via the weak interaction, thereby depleting the branching ratio for $J_T \to H^0 + \gamma$. Since the

J_T production cross section both in e^+e^- annihilation and pp and $p\bar{p}$ collisions decreases with increasing m_{J_T}, we feel that an optimal situation for Higgs search is in the range $m_{J_T} \simeq (40\text{-}60)$ GeV and of course with $m_{H^0} < m_{J_T}$.

The signatures for $J_T \to H^0 + \gamma$ have been worked out in detail in Ref. 21, where it was shown that one could search for peaks in the inclusive photon energy distribution and analyze the hadronic junk recoiling against the photon. Since $BR(H^0 \underset{b\bar{b}}{\to} l l^{\mp} + X) \simeq 0.64$, the hadronic jet will very <u>often contain a charged lepton</u> which would provide discrimination from decays of the type

$$J_T \to \gamma + GG \tag{4.25}$$
$$\qquad\qquad \hookrightarrow \text{glueball} \to \text{hadrons}$$

which otherwise could mock the decay

$$J_T \to \gamma + H^0 \tag{4.26}$$
$$\qquad\qquad \hookrightarrow \text{hadrons}$$

In Table 1, we show the relative branching ratio $R_{H^0\gamma/\mu^+\mu^-} \equiv \Gamma(J_T \to H^0\gamma)/\Gamma(J_T \to \mu^+\mu^-)$ for some representative values of m_{J_T} and m_{H^0}.

The production of toponium in both the e^+e^- annihilation and pp ($p\bar{p}$) collisions has acquired a new interest, namely that it provides one of the most hopeful reactions to discover the standard Weinberg-Salam Higgs if the phase space permits the decay (4.26). Perhaps it is worth remarking that in theories with dynamical symmetry breaking, the branching ratio for the process involving a color singlet neutral PGB, π'^0

$$J_T \to \pi'^0 + \gamma \tag{4.27}$$

is comparable to the decay rate for $J_T \to H^0 + \gamma$. We quote the relative rates[22]

$$\frac{\Gamma(J_T \to \pi'^0 + \gamma)}{\Gamma(J_T \to H^0 + \gamma)} \approx (n_{F'})^{-1} \tag{4.28}$$

where $n_{F'}$ is the number of hyperquark doublets. The dominant decays of $O(10\ \text{GeV})$ π'^0 are similar to those of H^0, namely

$$\pi'^0 \to b\bar{b},\ c\bar{c},\ \tau^+\tau^-$$

though the details may differ.[23] We shall explore the search of the radiative decay $J_T \to H^0 + \gamma$ in pp and $p\bar{p}$ collisions in the next section.

ii. HIGGS PRODUCTION IN pp AND p$\bar{p}$ COLLISIONS

DIRECT H^o PRODUCTION

The subprocesses that could lead to direct H^o production in pp and p$\bar{p}$ collisions are shown in Fig. 7. Thus, one could have a Drell-Yan type production

$$q + \bar{q} \to H^o \qquad (4.29)$$

as well as the gluon fusion mechanism of Ref. 24

$$G + G \to H^o \qquad (4.30)$$

leading to

$$p + p(\bar{p}) \to H^o + X \qquad (4.31)$$

The production cross section for (4.31) through the subprocess (4.29) can be expressed in terms of rapidity distribution

$$\frac{d\sigma_H}{dy} \sim \frac{\pi}{6} \Sigma g_q^2 \, \tau m_H^{-2} [F_q(\tau^{1/2} e^y) F_{\bar{q}}(\tau^{1/2} e^{-y}) + q \leftrightarrow \bar{q}] \qquad (4.32)$$

where

$$g_q = (\sqrt{2}\, G_F)^{1/2} m_q$$
$$y = \frac{1}{2} \ln \left(\frac{E + p}{E - p}\right) \; ; \; \tau = m_H^2/s \qquad (4.33)$$

$\sqrt{s}$ is the center-of-mass energy and $F_q(F_{\bar{q}})$ stands for the distribution function of finding a quark of flavor $q(\bar{q})$ inside the hadron. It should be remarked that m_q in the Yukawa coupling g_q is the current algebra mass because this is what the Higgs mechanism generates. Thus the coupling g_q is sizeable only for the heavy quarks but alas there are no heavy quarks in the proton! One could excite them from the proton (antiproton), but even at Q^2 available at ISABELLE and TEVATRON energies one would not excite enough $c\bar{c}$, $b\bar{b}$ and $t\bar{t}$ to make (4.32) appreciable. So, I shall neglect the subprocess (4.29).

Let us now consider the gluon fusion mechanism, Eq. (4.30). This leads to a rapidity distribution[24]

$$\frac{d\sigma_H}{dy} \sim \frac{\pi}{32} \left[\frac{\alpha_s(m_H^2)}{\pi}\right]^2 \frac{G_F}{\sqrt{2}} \frac{N_F^2}{9} \tau \, F_G(\tau^{1/2} e^y) \, F_G(\tau^{1/2} e^{-y}) \qquad (4.34)$$

where $\alpha_s(m_H^2)$ is the QCD running coupling constant evaluated at $Q^2 = m_H^2$ and N_F is the number of heavy flavors ($m_Q > 0.2\, m_H$). The process (4.34) is small due to $\alpha_s(m_H^2)^2$. Using the experimental result

$$\int_0^1 \xi \, F_G(\xi) \, d\xi = 1/2 \qquad (4.35)$$

and using $\xi F_G(\xi) = 3(1 - \xi)^5$, $N_F = 3$, one gets

$$\sigma(pp \to GG \to H^o + X) = \sigma(p\bar{p} \to GG \to H^o + X) \simeq 50 \text{ Pb} \qquad (4.36)$$
$$\text{for } \sqrt{s} = 800 \text{ GeV}; \; m_H = 11 \text{ GeV}.$$

The cross section for the gluon fusion mechanism is shown in Fig. 8 as a function of $\sqrt{s}$. Note that the cross section for (4.31) at the ISABELLE and TEVATRON energies is not small.

To have some idea about the signal to background ratio, let us calculate the cross section $\sigma(pp \to H^o + X)$, $H^o \to \mu^+\mu^-$, for the optimistic case of $m_{H^o} \simeq 10$ GeV, with $BR(H^o \to \mu^+\mu^-) \simeq 2 \times 10^{-3}$. This gives for $\sqrt{s} \simeq 800$ GeV

$$\sigma(pp \to H^o + X) \simeq 10^{-37} \text{ cm}^2, \quad H^o \to \mu^+\mu^- \qquad (4.37)$$

The Drell-Yan background evaluated for $m(\mu^+\mu^-) \simeq 10$ GeV at $\sqrt{s} \simeq 800$ GeV is[25]

$$\sigma(pp \underset{DY}{\to} \mu^+\mu^- + X) \simeq 10^{-34} \text{ cm}^2 \qquad (4.38)$$

Thus, the $\mu^+\mu^-$ mode looks hopeless. It is clear that one has to look for some other trigger on H^o. I have no wisdom to offer except urging my experimental colleagues to start thinking about developing techniques to tag the $\tau^\pm$ produced in hadron-hadron collisions and develop jet mass reconstruction techniques,[26] which might ultimately help establish the Higgs signal. For pure fun let me quote the cross section (for $m_{H^o} \simeq 10$ GeV)

$$\sigma(pp \to H^o + X) \simeq O(10^{-34} \text{ cm}^2), \quad H^o \to \tau^+\tau^- \qquad (4.39)$$

which is comparable to the Drell-Yan process at $\sqrt{s}$ = 800 GeV for $m(\mu^+\mu^-) \simeq$ 10 GeV.

$$\sigma(pp \to \tau^+\tau^- + X) \simeq O(10^{-34}\ \text{cm}^2) \qquad (4.40)$$

ASSOCIATED H^o PRODUCTION

The associated production of H^o in pp and $p\bar{p}$ collisions goes via the production of Z^o

$$pp(\bar{p}) \to Z^o + X \qquad (4.41)$$

followed by the decays of Z^o

$$Z^o \to H^o + \gamma$$

and

$$Z^o \to H^o + Z^o, \quad Z^o \to \mu^+\mu^-$$

discussed in the last section.* I am afraid that the cross sections involved in any useful tag are of order 10^{-39} cm^2 or less at ISABELLE energies,[18] making the associated H^o production mechanisms unattractive for the Higgs search.

PRODUCTION OF H^o THROUGH H^o-P_b MIXING

This mechanism has been proposed by Ellis et al.[19] The point is that if the Coleman-Weinberg estimate of m_{H^o} is right, then we expect m_{H^o} to be in the mass range 10-11 GeV, though the one-loop result prefers $m_{H^o} \approx$ 11 GeV. Thus if m_{H^o} is close in mass to the $P_b(\equiv {}^3P_o)$ state of the Υ-family, then there could be appreciable H^o-P_b mixing leading to the process[19,27]

$$pp(\bar{p}) \to P_b + X, \quad P_b \to H^o \qquad (4.42)$$

Ruckl and Baier (see these proceedings) have recently calculated the inclusive production of the P_b state in pp and $p\bar{p}$ collisions and it is substantial. The production of H^o through (4.42) then depends on the mixing parameter C, which can be calculated in a non-relativistic quark model calculation[19]

$$C \simeq \left[\frac{27\sqrt{2}}{\pi} G_F\, m_{P_b}\, |R'(0)|^2\right]^{1/2} \qquad (4.43)$$

where R'(0) is the derivative of the wave-function at the origin. The decay mode which is most favorable is

*) In addition one has the associated production

$$pp(\bar{p}) \to Z^o_{\text{vir.}} + X$$
$$\to Z^o + H^o$$

$$P_b \underset{\text{Mixing}}{\to} H^o \to (\tau^+\tau^-,\ c\bar{c}) \qquad (4.44)$$

Putting everything together, one gets

$$\frac{\Gamma(P_b \to \tau^+\tau^-)}{\Gamma(P_b \to \text{all})} \simeq \frac{9\, G_F^2\, m_\tau^2\, m_{P_b}^6 \left(1 - \frac{4m_\tau^2}{m_{P_b}^2}\right)^{3/2}}{128\, \pi^2\, \alpha_s^2 (m_H^2)(m_H^2 - m_{P_b}^2)^2} \simeq \frac{2.5 \times 10^{-7}}{(m_H - m_{P_b})^2} \qquad (4.45)$$

where $\Delta H = m_H - m_{P_b}$ is expressed in GeV. For ΔM = 20 MeV, (4.45) gives

$$\frac{\Gamma(P_b \to \tau^+\tau^-)}{\Gamma(P_b \to \text{all})} \simeq 6 \times 10^{-4} \qquad (4.46)$$

which is rather small and presumably realistic if m_{H^o} is close to m_{P_b}. However, as an extreme example one could think of a scenario in which $\Delta M <$ 1 MeV such that $\Delta M \sim \Gamma(P_b \to \text{all})$. In that case there will be complete mixing in the P_b-H^o sector, making the perturbation theory estimate (4.43) of the mixing parameter inapplicable. In the event of complete mixing, one would have a situation very similar to the K^o-$\bar{K}^o$ system namely, there will be two states of P_b with

$$P_{b\binom{1}{2}} = \frac{1}{\sqrt{2}} (|H\rangle \pm |P_b\rangle) \qquad (4.47)$$

having branching ratio $BR(P_b \to \tau^+\tau^-) \sim 10\%$ for both the states P_{b_1} and P_{b_2}! The prediction (4.47) has not yet been checked but it is clear that its validity would require an accidental degeneracy of m_{H^o} and m_{P_b} and in general such accidents are not very widespread in nature!

Perhaps, it is fair to say that if $\Delta M \simeq$ 100 MeV, then the scenario (4.42) would lead to a marginal increase in the direct H^o production at the ISABELLE and TEVATRON energies.

PRODUCTION OF H^o THROUGH THE RADIATIVE DECAY $J_T \to H^o + \gamma$

In the last section I have discussed the exciting possibility of the process

$$J_T \to H^o + \gamma$$

which could be a large fraction of the rate for $J_T \to \mu^+\mu^-$, the traditional way of hunting for the vector mesons in hadron-hadron collisions. In this section I would like to advocate using the radiative decay of J_T to search for H^o in the pp and $p\bar{p}$ collisions at ISABELLE and TEVATRON. The process one should search for is

$$pp(\bar{p}) \to J_T + X, \quad J_T \to H^o + \gamma, \quad H^o \to 1\ell^\pm + X' \tag{4.48}$$

Depending on the mass difference $m_{J_T} - m_{H^o}$, the photon in (4.48) will be very energetic and will have a large p_T recoiling against a jet whose composition will depend on the mass of the H^o. If $m_{H^o} < 2m_B$, then the decay mode

$$H^o \to \tau^+\tau^-$$

could be as big as 25% of all the H^o decays, and one could look for

$$pp(\bar{p}) \to J_T + X, \quad J_T \to H^o + \gamma\,(\text{large } p_T), \quad H^o \to \tau^+\tau^- \tag{4.49}$$

On the other hand if $m_{H^o} > 2m_B$, then the decays of H^o would be dominated by

$$H^o \to b\bar{b}$$

leading to the final state in (4.48). Since the branching ratio $H^o \to b\bar{b} \to 1\ell^\pm + X$ is about 2/3, there won't be any appreciable loss in the event rate. The requirement of large-p_T photon would reduce the background to (4.49) from the usual Drell-Yan background $pp(\bar{p}) \to J_T + X$, $J_T \to \tau^+\tau^-$ and $pp(\bar{p}) \to \tau^+\tau^- + X$.

Requiring a prompt lepton in the hadronic shower recoiling against the large-p_T photon would reduce the background from the inclusive photon production background

$$pp \to \gamma + X \tag{4.50}$$

Next, we would like to calculate the rate for the process $pp \to J_T + X$, $J_T \to H^o + \gamma$ and compare it with the rate for the background (4.50). There is a considerable amount of uncertainty involved in estimating the toponium production rate in pp and $p\bar{p}$ collisions. This is so because for large values of m_t, the wave-function becomes coulombic. In addition there are uncertainties about the diffractive component of $pp(\bar{p}) \to J_T + X$ as well as the contribution of the χ_T states to J_T production via the reaction $pp(\bar{p}) \to \chi_T + X$, $\chi_T \to J_T + \gamma$. To make a ball park estimate, we shall use Gaisser scaling,[28] which works for the J/ψ and Υ production in pp collisions within a factor 2. The quantity which scales is

$$\frac{d\sigma}{dy} = \aleph^o\, \Gamma_{3g}\, \frac{d\sigma^{DY}}{dMdy} \tag{4.51}$$

where the coefficient $\aleph^o$ has been calculated to be $\aleph^o = 15 \times 10^6$.[28] Using this estimate for the production cross section, the rate for the quantity

$$(pp \to J_T + X)\ BR(J_T \to H^o + \gamma) \tag{4.52}$$

is tabulated in Table 2 for various representative values of m_{J_T} and m_{H^o} = 11 GeV at the ISABELLE energy $\sqrt{s}$ = 720 GeV. The numbers in the last column are the expected number of events assuming an integrated luminosity of 3.6×10^{39} (corresponding to 1000 hours of running time with $\mathcal{L} = 10^{33}$). We note that the event rates are quite encouraging. Rates for other allowed values of m_{H^o} can be obtained by combining Tables 1 and 2. For $m_{J_T} > m_Z$, the branching ratio $J_T \to \mu^+\mu^-$ becomes very small.

Let us now calculate the cross section for the background

$$pp(\bar{p}) \to \gamma + X$$

where the photon is produced, for example, in the subprocess $G + q \to \gamma + q$. The differential cross section $d\sigma/dp_T d\Omega$ has been evaluated in Ref. 25 to be

$$\left.\frac{d\sigma}{dp_T d\Omega}\right|_{\gamma_{p_T} = 30\text{ GeV}} \simeq 1\text{ Pb} \tag{4.53}$$

which is comparable to the production rate for the process

$$pp \to J_T + X, \quad J_T \to H^o + \gamma \qquad (4.54)$$

However, the background can be very much reduced by triggering in addition on a prompt lepton. One could even determine m_{H^o} by constructing the energy and momentum of the hadronic shower recoiling against the photon and measuring $d\sigma/dm_{\gamma "H^o\ jet"}$. This needs some effort and experience with hadronic jets at ISABELLE and TEVATRON energies.

The prospects of observing a Higgs signal in pp and $p\bar{p}$ collisions at high energies through the process (4.54) at the level of O(1 Pb) is exciting. One should point out that direct H^o production in pp and $p\bar{p}$ collisions through the gluon fusion mechanism is quite large at ISABELLE and TEVATRON energies. However, the Higgs so produced lacks a reliable trigger. Thus, the somewhat smaller production cross section for (4.54) is compensated by the relatively clean trigger that we have here advocated.

iii. HIGGS PRODUCTION IN LEPTON-NUCLEON PROCESSES

These processes include the reactions

$$\begin{aligned} \ell + N &\to \nu_\ell + H^o + X \\ &\to \ell + H^o + X \\ \nu_\ell + N &\to \nu_\ell + H^o + X \\ &\to \ell + H^o + X \end{aligned} \qquad (4.55)$$

The lowest order processes involving the bremstrahlung off the lepton line are negligible. The next order diagrams involving double $W^\pm$ and Z^o exchanges are shown in Fig. 9. One could express the differential distribution corresponding to these diagrams as[29]

$$\frac{d^4\sigma}{dxdydx'dy'} = \frac{G_F^3 s^2 x'}{4\sqrt{2}\pi^3}\left[(1 + \frac{s}{m_W^2}xy)^2(1 + \frac{s}{m_W^2}x'y')^2\right]^{-1}$$

$$\times [\{u(x',Q'^2) + c(x', Q'^2)\}f_1$$

$$+ \{\bar{d}(x',Q'^2) + \bar{s}(x', Q'^2)\}f_2] \qquad (4.56)$$

with $Q'^2 = -q'^2 = sx'y'$,

$$x = \frac{-q^2}{2p\cdot q},\ x' = \frac{-q'^2}{2p\cdot q'},\ y = \frac{q\cdot p}{k\cdot p},\ y' = \frac{q'\cdot p}{k\cdot p}$$

f_1 and f_2 are kinematic factors for collisions of fermions with like and unlike helicities and are given by

$$f_1 = 1 - y + y' - \frac{x}{x'} - \frac{y'}{y} + \frac{xy}{x'} + \frac{2xy'}{x'y} - \frac{m_H^2}{s}\frac{(1-y)}{x'y}$$

$$f_2 = (1 - y)\left[1 - \frac{x}{x'} - \frac{y'}{y} - \frac{xy'}{x'} + \frac{2xy'}{x'y} - \frac{m_H^2}{s}\frac{1}{x'y}\right] \qquad (4.57)$$

The functions $u(x', Q'^2)$ etc., are the quark densities inside nucleon corresponding to the u quarks. In the low energy limit (i.e. $m_W \to \infty$), and neglecting any Q^2 dependence of the structure functions, one has

$$\sigma \sim \frac{G_F^3 s^2}{8\sqrt{2}\pi^3}\int_0^1 dx\,\frac{1}{6}\,x\,F_2(x) \qquad (4.58)$$

using $\sigma(\nu_\mu \to \mu^-) = G_F^2 s/\pi$ gives

$$\frac{\sigma(\nu_\mu \to \mu^- + H^o + X)}{\sigma(\nu_\mu \to \mu^- + X)} \sim 3 \times 10^{-8}\left(\frac{E_\nu}{M_N}\right) \qquad (4.59)$$

The cross section for other processes in (4.55) are even smaller. With $\sin^2\theta_W \simeq 0.22$, one has

$$\sigma(\nu \to \nu_H) = 2\sigma(\bar{\nu} \to \bar{\nu}_H) \simeq 0.04\,\sigma(\nu_\mu \to \mu^- + H) \qquad (4.60)$$

and it looks hopeless.

At high energies, the results are shown in Fig. 10 for both the cases involving Q^2-independent and Q^2-dependent quark densities. Typically

$$\begin{aligned} \sigma(e + p \to \nu_e + H^o + X) &< 10^{-37}\ cm^2 \quad \text{for } s = 10^7\ GeV^2 \\ &< 10^{-39}\ cm^2 \quad \text{for } s = 10^4\ GeV^2 \end{aligned} \qquad (4.61)$$

with QCD corrections reducing the rate by roughly a factor 2.

My conclusion is that ep and νp machines are not suitable to study the question of Higgs production at high energies.

iv. HIGGS PRODUCTION IN e^+e^- ANNIHILATION

e^+e^- machines are probably the last hope of the broken hearts who would like to see a Higgs in their life! The history of e^+e^- machines has been a success story which primarily could be traced back to their role as vector boson factories. Depending upon the energy resolution and the mass of the resonance as well as its width, the enhancement at these vector meson poles could be several thousand units of R, which compensates the intrinsically smaller cross section of the process $e^+e^- \to$ hadrons. Of course, the other advantage is that the initial state is very precisely known.

We have already discussed the potential role that vector bosons, J_T and Z^o could play in Higgs searches through the decays $J_T \to H^o + \gamma_{monoenergetic}$ and $Z \to H^o + \mu^+\mu^-$. If one could produce these vector bosons in abundance, then the vector boson factories have a much cleaner environment to study these decays. However if $m_{H^o} > m_{Z/2}$ and if $m_{H^o} > m_{J_T}$ then the vector boson factories lose their edge on other machines as far as Higgs search is concerned.

I shall now discuss some other processes in the context of e^+e^- machines.

1. $e^+e^- \xrightarrow{\gamma,Z} H^oH^o$

This is not allowed due to Bose symmetry.

2. Bremsstrahlung mechanisms

Here a Higgs can be bremsgestrahlt off the initial or final fermion line[2,27]

$$\frac{\sigma(e^+e^- \to \mu^+\mu^-H^o)}{\sigma(e^+e^- \to \mu^+\mu^-)} \sim G_F\, m_e^2\, \ell ns \simeq 2.2 \times 10^{-11} \qquad (4.62)$$

$$\frac{\sigma(e^+e^- \to \tau^+\tau^-H^o)}{\sigma(e^+e^- \to \mu^+\mu^-)} \sim 2.8 \times 10^{-4} \qquad (4.63)$$

$$\frac{\sigma(e^+e^- \to t\bar{t}H^o)}{\sigma(\mu^+\mu^-)} \geq 0.1 \quad \text{for } m_t \geq 20 \text{ GeV} \qquad (4.64)$$

However, it is very hard to identify H^o in the final state $t\bar{t}H^o$ since it lacks a good trigger.

3. $e^+e^- \to H^o$

$$\sigma(e^+e^- \to H^o) \simeq \frac{4\pi}{m_H^2} \frac{\Gamma(H^o \to e^+e^-)}{\Gamma(H^o \to \text{all})} \qquad (4.65)$$

assuming

$$\Gamma(H^o \to Q\bar{Q}) \approx \Gamma_{total}(H^o)$$

$$\to BR(H^o \to e^+e^-) \simeq \frac{1}{3}\left(\frac{m_e}{m_Q}\right)^2$$

which gives

$$\begin{aligned} \Delta R(e^+e^- \to H^o) &\simeq \frac{3}{\alpha^2} BR(H^o \to e^+e^-) \\ &\simeq \frac{1}{\alpha^2} (m_e/m_Q)^2 \qquad (4.66) \\ &\simeq 2 \times 10^{-3} \text{ for } 2m_D < m_{H^o} < 2m_B \\ &\simeq 2 \times 10^{-4} \text{ for } 2m_B < m_{H^o} < 2m_T \end{aligned}$$

4. $e^+e^- \to H^o + \gamma$ [30]

The one-loop diagrams are shown in Fig. 11 and one could express the result as

$$\frac{\sigma(H^o\gamma)}{\sigma(\mu^+\mu^-)} \simeq 4.8 \times 10^{-10}\, s(1 - m_H^2/s)^3 \left|I_W - \sum_F I_F\right|^2 \qquad (4.67)$$

$$I_W \Big|_{\sqrt{s} = 30 \text{ GeV}} \simeq 1.7$$

$$I_F \Big|_{m_F \to \infty} \simeq 1/3 \text{ for } \binom{\nu_L}{L^-} \text{ doublet}$$

$$\simeq 5/9 \text{ for } \binom{G}{H} \text{ doublet}$$

This gives

$$\Delta R(H^o\gamma) \simeq O(10^{-6}) \quad (4.68)$$

$$\text{for } m_{H^o} \lesssim 20 \text{ GeV}$$
$$20 \text{ GeV} \lesssim \sqrt{s} \lesssim 90 \text{ GeV}$$

an even difficult proposition than measuring $\Delta R(e^+e^- \to H^o)$!

5. $e^+e^- \to Z^oH^o$ [31]

This is the last process that I would like to discuss in my talk. Since the Z^oZ^oH coupling is large, one expects a huge cross section as the center-of-mass energy in e^+e^- annihilation increases. The cross section can be expressed as

$$R_{ZH^o} = \frac{\sigma(Z^oH^o)}{\sigma(\mu^+\mu^-)} = \frac{3}{64}\left(\frac{m_Z}{38 \text{ GeV}}\right)^4 \frac{m_Z}{2\sqrt{2}\, m_H} \times [1 + (1 - 4\sin^2\theta_W)^2] \quad (4.69)$$

which is not a small number. The ratio R_{ZH^o} peaks at $\sqrt{s} = m_Z + \sqrt{2}\, m_H$ and the rates at the peak assuming various values of m_{H^o} are shown in Table 3. The Z^o could be identified through

$$Z^o \to e^+e^-,\ \mu^+\mu^-$$

then

$$m_H^2 = (\sqrt{s} - E_Z)^2 - p_Z^2 \quad (4.70)$$

The event rates corresponding to

$$e^+e^- \to Z^oH^o, \quad Z^o \to \ell^+\ell^- \quad (4.71)$$

are shown in the last column assuming optimistically a luminosity of 10^{31}. The table serves to show that LEP could at most see a Higgs of $m_{H^o} \lesssim m_{Z/2}$. The situation should be contrasted with the Higgs production through the gluon fusion mechanism in pp and $p\bar{p}$ collisions. My feeling is that at ISABELLE and TEVATRON, one would be able to explore a much larger range of m_{H^o} --- provided one could trigger on $\tau^+\tau^-$ and/or on the energetic jets in the decays of H^o.

V. CONCLUSION

The problem of understanding the mechanism that breaks chiral symmetry and gives fermions and gauge bosons their masses is not yet understood. Higgs mechanism is a possible solution. The standard Weinberg-Salam theory has a neutral Higgs, which is a physical particle. It's detection will strengthen our belief in the underlying framework. I reviewed attempts made to harness the Higgs.

Among the conceivable production mechanisms in proton-proton and proton-antiproton collisions, the most hopeful place is the process

$$pp(\bar{p}) \to J_T + X, \quad J_T \to H^o + \gamma \quad \text{if } m_{H^o} < m_{J_T}$$

since it provides a rather powerful trigger. It will however need a high luminosity pp or $p\bar{p}$ collider, for example the phase II of ISABELLE. The gluon fusion mechanism $pp(\bar{p}) \to GG \to H^o + X$ has a potentially large cross section. It is imperative to develop a reliable trigger on H^o --- other than $\mu^+\mu^-$.

The production of vector boson J_T and Z^o in e^+e^- annihilation and the decays $J_T \to H^o + \gamma$, $Z \to H^o + \mu^+\mu^-$ and the reaction $e^+e^- \to Z^oH^o$ are the only other hopes of finding an H^o. However if m_{H^o} is large [$\gtrsim$ O(100 GeV)], then e^+e^- machines won't be able to see them simply because these machines run out of gas beyond $\sqrt{s} = 200$ GeV. In that event pp and $p\bar{p}$ machines are the only hopes to see a Higgs, or any other particle or phenomenon, replacing the Higgs mechanism. Only pp and $p\bar{p}$ machines provide energies reaching $\sqrt{s} = \langle\phi^o\rangle \sim 250$ GeV in a channel which could communicate with the vacuum. I am sure that when the pp and $p\bar{p}$ machines will probe the electroweak vacuum at $\sqrt{s} \simeq \langle\phi^o\rangle$, we are destined to observe new and fascinating phenomenon.

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TABLE 1

The ratio $\frac{\Gamma(J_T \to H^o + \gamma)}{\Gamma(J_T \to \mu^+\mu^-)}$ in the Weinberg-Salam model

m_{H^o} (GeV) \ $m(J_T)$ (GeV)	40	60	80
10	.14	.32	.58
20	.11	.29	.55
30	.065	.24	.50
40		.18	.44
50		.10	.35
60			.25
70			.13

Table 2

Production cross section for the process $pp \to J_T + X \to H^o + \gamma + X$ using Gaisser scaling. The numbers correspond to using $m_{H^o} \simeq 11$ GeV, $\sqrt{s} = 720$ GeV, $BR(H^o \to 1\ell^\pm + X) = 0.64$ and an integrated luminosity of 3.6×10^{39}.

$M(J_T)$ (GeV)	$\sigma(pp \to J_T + X)$, $J_T \to H^o + \gamma$	# Events $pp \to J_T \to H^o + \gamma$, $\to \geq 1\ell^\pm + X$
40	1.0 Pb	2300
50	0.5 Pb	1150
60	0.22 Pb	500
70	0.1 Pb	230
80	0.04 Pb	∿ 100

TABLE 3

Expected rates for the process $e^+e^- \to Z^oH^o$ at the peak value $\sqrt{s} = m_Z + \sqrt{2}\, m_H$. The entries in the last two columns are obtained for $\mathcal{L} = 10^{31}$ cm^{-2} sec^{-1} and a branching ratio $Z^o \to \ell^+\ell^- = 3\%$ (Barbiellini et al. in Ref. 2).

M_H(GeV)	$\sqrt{s}$ (GeV)	$\frac{\sigma(Z^o + H^o)}{\sigma_{pt}}$	$\frac{\# H^o + Z^o}{10^3 \text{ hrs.}}$	$\frac{\# H^o\ell^+\ell^-}{10^3 \text{ hrs.}}$
10	104	4.7	1458	44
30	132	1.56	280	∿ 9
45	154	1.04	138	∿ 4
60	175	0.78	82	∿ 3

FIGURE CAPTIONS

Fig. 1 Feynman diagrams for the decay $H^o \to 2\gamma$.

Fig. 2 Feynman diagram for the decay $H^o \to 2G$.

Fig. 3 Feynman diagram forthe decay $Z^o \to H^o + \gamma$.

Fig. 4 Decay rates for $\Gamma(Z^o \to H^o + \gamma)/\Gamma(Z^o \to \mu^+\mu^-)$ and $\Gamma(Z^o \to H^o + \mu^+\mu^-)/\Gamma(Z^o \to \mu^+\mu^-)$ for different values of m_H/m_Z from Ref. 15.

Fig. 5 Dimuon invariant mass distribution in the decay $Z^o \to H^o + \ell^+\ell^-$ from Ref. 16.

Fig. 6 Distribution of the dilepton invariant mass $m(\ell^+\ell^-)$ in $pp \to \ell^+\ell^-\ell^\pm X$ at $\sqrt{s} = 800$ GeV from the sources

(i) decays of $b\bar{b}$ and $t\bar{t}$ produced in pp collisions.

(ii) Drell-Yan process $pp \to (\gamma, Z^o) \to \ell^+\ell^-$ along with a third lepton produced with a probability 10^{-4} and

(iii) $pp \to Z^o + H^o + X$ followed by $Z^o \to \ell^+\ell^-$ and $H^o \to \ell^\pm + X$.

Label $\ell_i\ell_j$ denotes leptons originated from particle i and j respectively and $(\ell^+\ell^-)_i$ denotes leptons from the same parent, from Ref. 18.

Fig. 7 (a) 'Drell-Yan' mechanism for $p + p(\bar{p}) \to H^o + X$.

(b) Gluon fusion mechanism for $p + p(\bar{p}) \to H^o + X$.

Fig. 8 Production cross section for $NN \to GG \to H^o + X$ ($N = p$ or $\bar{p}$) for $m_{H^o} = 11$ GeV (from W.Y. Keung).

Fig. 9 Feynman diagrams for Higgs production in lepton-nucleon scattering

processes $\ell + p \to \nu_\ell + H^o + X$

and $\ell + p \to \ell + H^o + X$.

Fig. 10 (a) Cross sections for $e + p \to \nu_e + H^o + X$ shown in Fig. 10 as a function of (CM energy)^2. Solid and dashed curves correspond to parton model with and without a QCD correction. Cases A, B and C are results for $m_{H^o} = 10$, 100 and 200 GeV.

(b) Same as (a) but for smaller range of s. Curves A, B and C are results with $m_{H^o} = 10$, 20 and 30 GeV, from Ref. 29.

Fig. 11 Diagrams for the process $e^+e^- \to H^o + \gamma$.

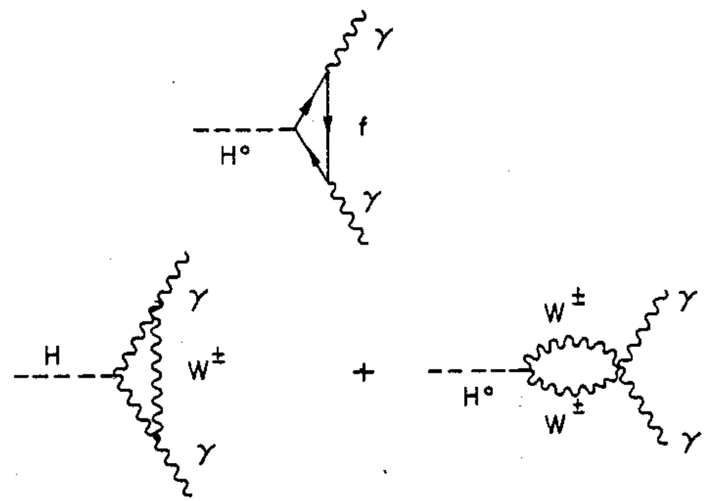

Fig. 1

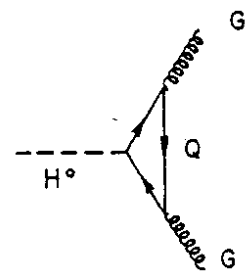

Fig. 2

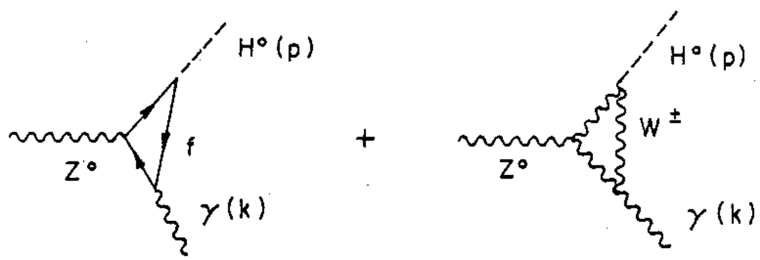
H° (p)
f
Z°
γ (k)
+
H° (p)
W ±
Z°
γ (k)

Fig. 3

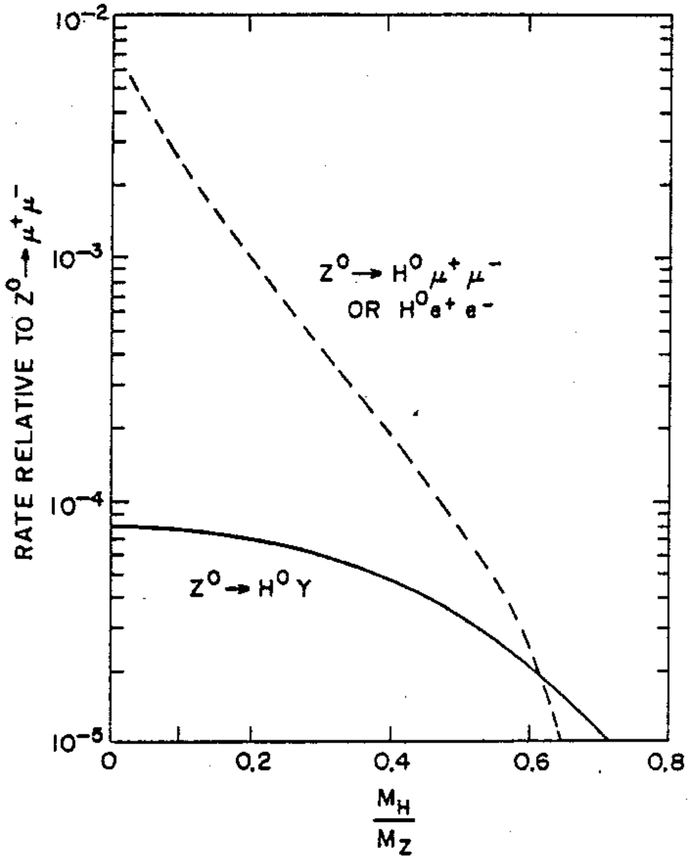
RATE RELATIVE TO Z⁰ → μ⁺μ⁻
10⁻²
10⁻³
10⁻⁴
10⁻⁵
Z⁰ → H⁰ μ⁺ μ⁻
OR H⁰e⁺ e⁻
Z⁰ → H⁰ γ
0
0.2
0.4
0.6
0.8
M_H / M_Z

Fig. 4

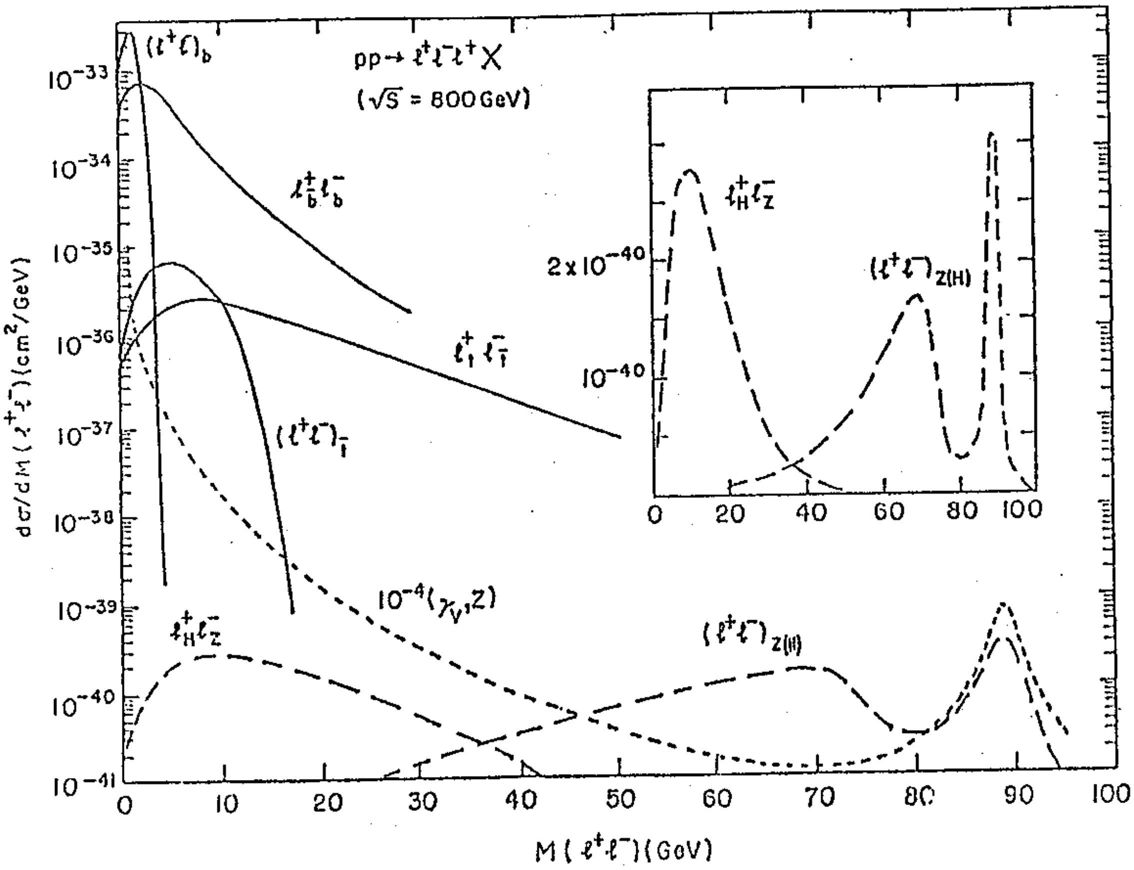
pp→ℓ⁺ℓ⁻ℓ⁺X
(√S = 800 GeV)
(ℓ⁺ℓ)b
ℓ⁺b ℓ⁻b
ℓ⁺1 ℓ⁻1
(ℓ⁺ℓ⁻)1
10⁻⁴(γV,Z)
ℓ⁺H ℓ⁻Z
(ℓ⁺ℓ⁻)Z(H)
dσ/dM (ℓ⁺ℓ⁻)(cm²/GeV)
M (ℓ⁺ℓ⁻)(GeV)
2x10⁻⁴⁰
10⁻⁴⁰

Fig. 6

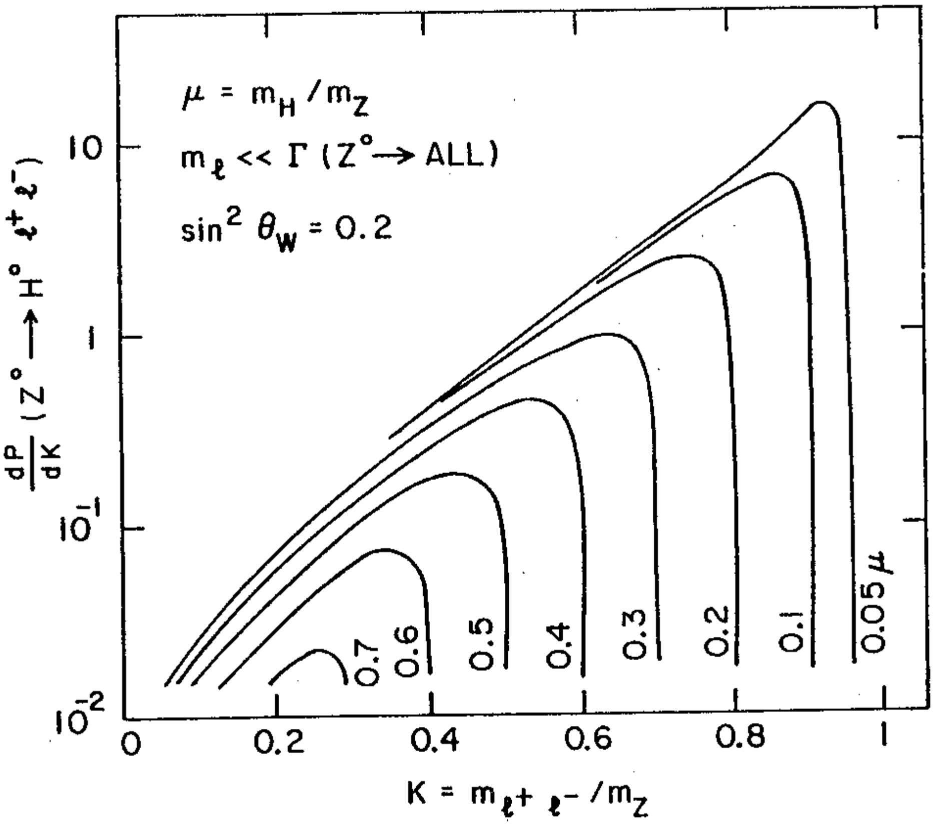
μ = mH / mZ
mℓ << Γ (Z°→ ALL)
sin² θW = 0.2
dP/dK (Z° → H° ℓ⁺ ℓ⁻)
K = mℓ+ ℓ- / mZ
0.7
0.6
0.5
0.4
0.3
0.2
0.1
0.05 μ

Fig. 5

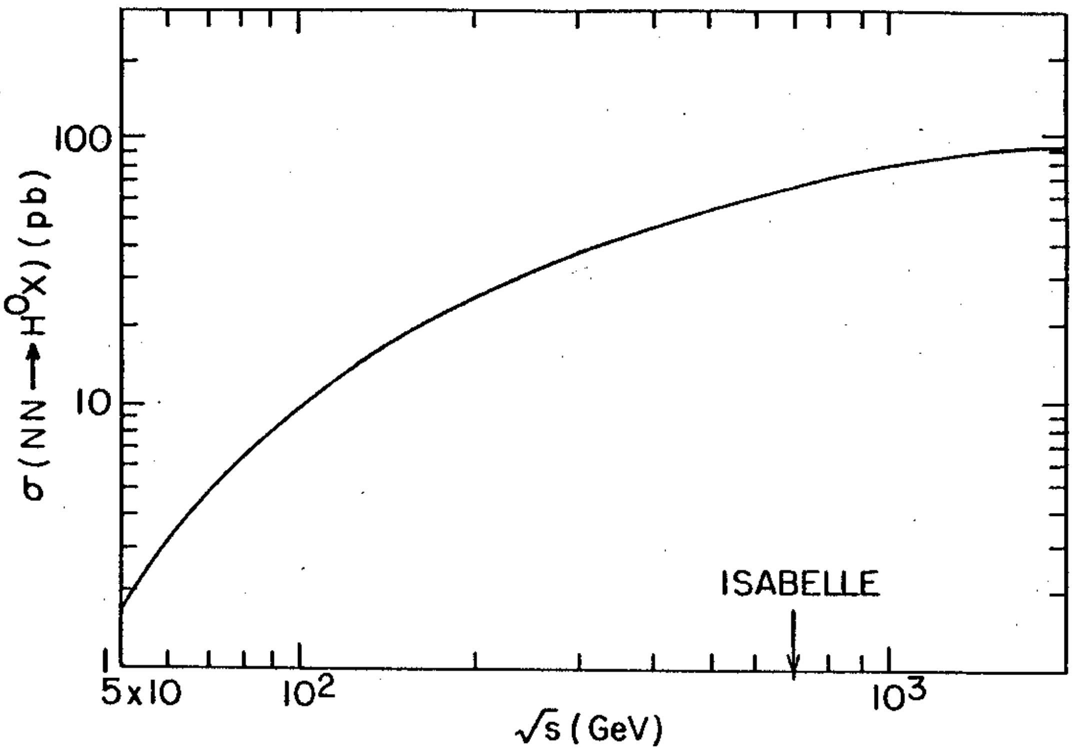

Fig. 8

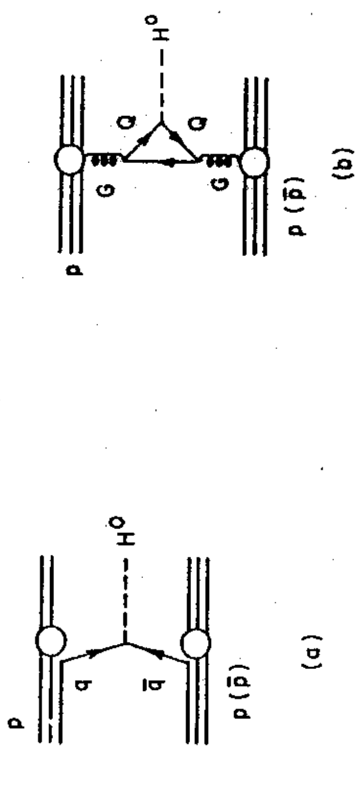

Fig. 7

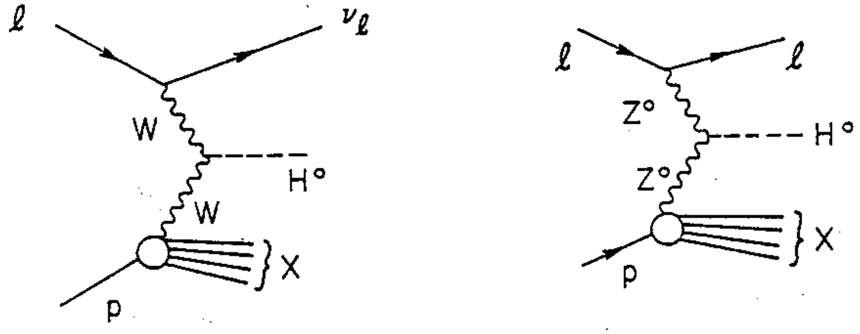
ℓ
ν_ℓ
W
H°
W
X
p
ℓ
ℓ
Z°
H°
Z°
X
p

Fig. 9

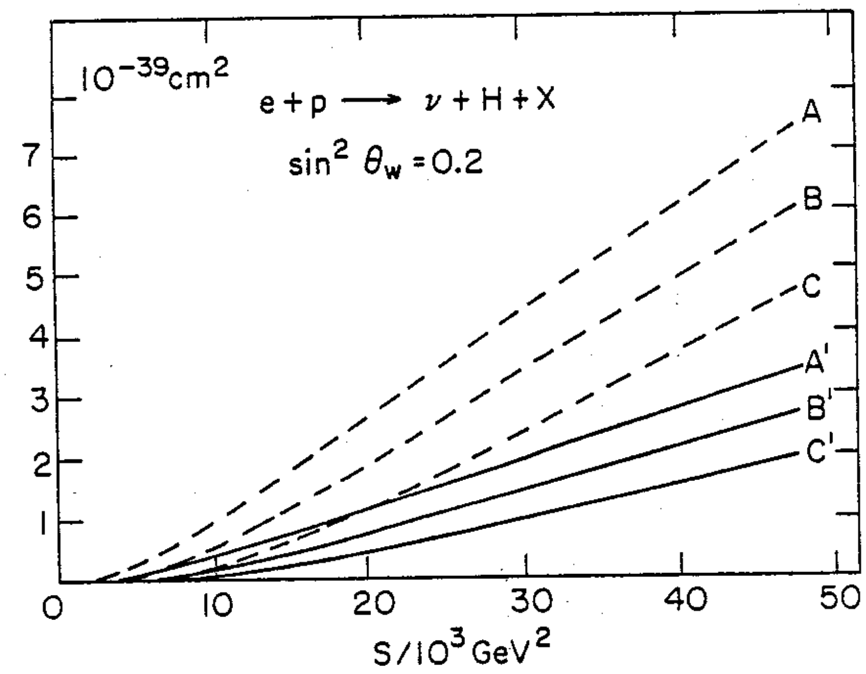
$10^{-39}cm^2$
e + p ⟶ ν + H + X
$\sin^2 \theta_W = 0.2$
A
B
C
A'
B'
C'
0
1
2
3
4
5
6
7
0
10
20
30
40
50
$S/10^3 GeV^2$

Fig. 10a

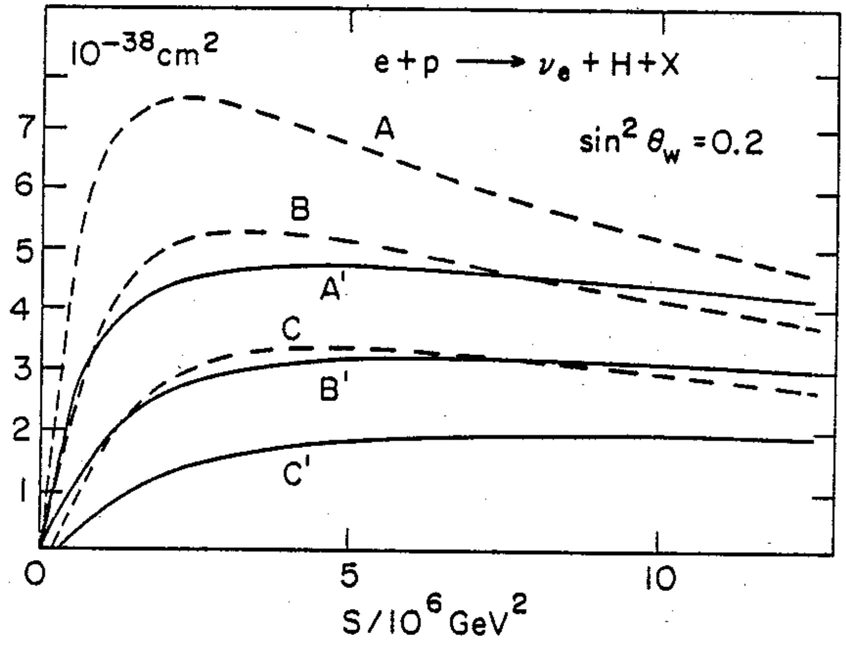

Fig. 10b

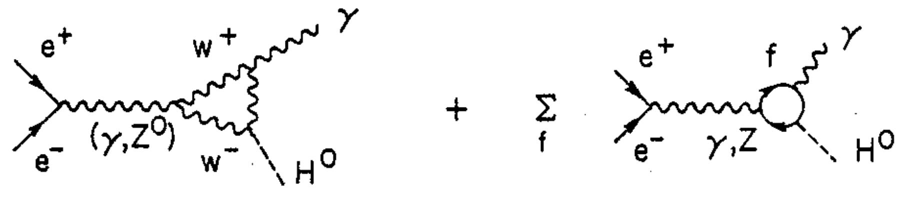

Fig. 11