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

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Abstract. The associated production of Higgs particles and Z (W) bosons in pp and $p\bar{p}$ collisions is studied. The total cross-section of the process $pp\,(p\bar{p}) \to HXl^+ l^-$ is estimated.

Keywords. Higgs bosons; Weinberg-Salam theory; $Z\,(W)$ boson decay; Drell-Yan approximation; quarks.

1. Introduction

At present it seems that all interactions between elementary particles can be described with the help of Gauge theories. The discovery of neutral currents and many successes of quantum chromodynamics (QCD) support this idea. QCD looks like the most suitable candidate for the theory of strong interactions. Moreover, the Gauge theories of weak and electromagnetic interactions are in good agreement with all the experimental data available now. Anyhow, the true test of these theories will come from the discovery of heavy vector bosons and those of Higgs bosons.

The Higgs mechanism of spontaneous symmetry breaking has now become an ingredient of unified theories of weak and electromagnetic interactions. The experimental discovery of Higgs bosons, thus, would emerge as a verification of one of the fundamental concepts of quantum electrodynamics (QED). Some complications occurring with Higgs bosons in Gauge theories, however, should be mentioned. This is first of all due to indeterminancy of Higgs sector in these theories. Being an optimist one can think that nature will choose the simplest possible version, *i.e.* Weinberg-Salam theory (Weinberg 1967; Salam 1968) with minimal Higgs sector (one Higgs boson). Further, the difficulty is that the mass of the Higgs boson (H-boson) is not predicted by the theory, only some poor limits have been indicated in the literature (Weinberg 1976; Linde 1976; Guth and Weinberg 1980).

An interesting mechanism of H-boson production in e^+e^- colliding beams ($e^+e^- \to ZH$) has been studied by several authors (Ellis *et al* 1976; Glashow *et al* 1978; Proeyen 1979; Sami and Fainberg 1981a). Glashow *et al* (1978) have proposed another mechanism for Higgs boson production.

$$pp \to X + Z\,(W^{\pm}) + H, \tag{1}$$

$$p\bar{p} \to X + Z\,(W^{\pm}) + H. \tag{2}$$

In fact, $Z(W^{\pm})$ bosons being unstable will decay soon into light generations, say leptons and quarks. Thus the study of processes (1) and (2) with the consideration of $Z(W^{\pm})$-boson decay, would have been more realistic. In this paper we have estimated the total cross-section of (1) and (2) with all possible decay of $Z(W^{\pm})$ boson into leptons, in the Weinberg-Salam theory. The production of light H-boson ($M_H \simeq 10$ GeV) in pp and $p\bar{p}$ collisions *via* high P_T bremsstrahlung from Z has been discussed by Finjord *et al* (1979).

2. Higgs bosons from fermion pair annihilation

For the sake of simplicity, we have done our calculations in simple parton model. According to this model, protons and antiprotons consist of quarks and antiquarks. Then the production of $Z(W^{\pm})$ bosons in (1) and (2) is interpreted as a result of an appropriate quark-antiquark pair annihilation. For this reason, we first of all study the elementary process (figure 1)

$$f\bar{f} \to Z(W^{\pm}) + H \to f'\bar{f}' + H, \tag{3}$$

where $f\bar{f}$, $f'\bar{f}'$ are the appropriate fermion pairs in initial and final states.

As in our previous paper (Sami and Fainberg 1981a), we use the covarient expression for $Z(W^{\pm})$ boson propagator instead of the usual Breit-Wigner formula:

$$[M^2 - Q^2 - \Pi(Q^2)]^{-1}, \tag{4}$$

where Q — four momentum of $Z(W^{\pm})$ boson,

$\Pi(Q^2)$—polarisation operator of unstable particle

$$M = M_Z, M_W.$$

From unitarity, the general expression for the total cross-section of all the processes going through $Z(W^{\pm})$ boson in $f\bar{f}$ colliding beams can be written as

$$\sigma_T \equiv \sigma_T(f\bar{f} \to) = J_\mu^{f\bar{f}}(p_+, p_-)\, 2\, I_m\, D^{\mu\nu}(Q^2)\, J_\nu^{f\bar{f}}(p_+, p_-) \Big/ \frac{Q^2}{2}$$

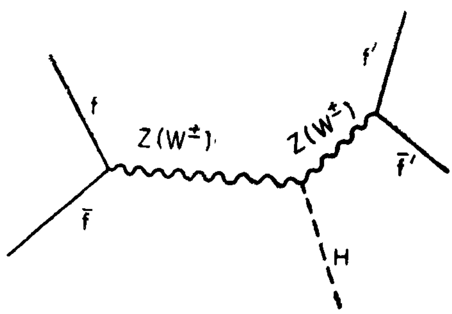


Figure 1. Graph for the process $f\bar{f} \to Z(W^{\pm}) + H \to f'\bar{f}' + H$, $f\bar{f}$, $f'\bar{f}'$ are the appropriate fermion pairs in initial and final states.

Here p_+, p_-, are the four-momenta of fermion and antifermion respectively, $Q^2 = (p_+, p_-)^2$ is the total energy squared of $f\bar{f}$ in the centre-of-mass system,

$$J_\mu^{f\bar{f}} = [C_L\,(\bar{\psi}_L\,\gamma_\mu\,\psi_L) + C_R\,(\bar{\psi}_R\,\gamma_\mu\,\psi_R)],$$

are the fermion currents in Weinberg-Salam theory and the constants C_L, C_R are given in table 1;

$I_m\,D_{\mu\nu}$ is the imaginary part of exact Z ($W^\pm$) boson propagator in our case.
Here and throughout, we have neglected the masses of fermions. If so, using (4) we have the expression for $I_m D_{\mu\nu}$

$$I_m D_{\mu\nu} = \frac{g_{\mu\nu}\,I_m\,\Pi(Q^2)}{[(M^2 - Q^2)^2 + (I_m\,\Pi(Q^2))^2]}, \tag{5}$$

where the real part of the polarization operator has been omitted. The whole contribution to $I_m\Pi$ comes from figures 2 and 3. Figure 2 corresponds to the process in which $Z(W^\pm)$ bosons produced from pair annihilation, at once decay into light generations. This process is more probable and starts before the appearance of Higgs boson. Figure 3 corresponds to the process in which $Z(W^\pm)$ bosons emit Higgs boson and then further decay into fermions (leptons and quarks).

$I_m\Pi(Q^2)$ can now be calculated as

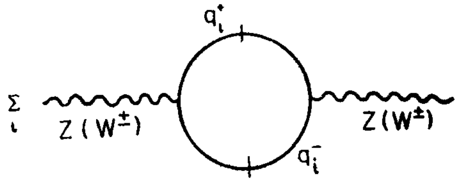
$$I_m\Pi(Q^2) = F_1 + F_2. \tag{6}$$

$$F_1 = \sum_i\,[(C_L^i)^2 + (C_R^i)^2]\;Q^2/24\pi \equiv M^2\,F_1\,(a),$$

Table 1. Constants C_L and C_R for all fermion pairs (quarks and leptons) interacting with $Z(W)$ bosons.

	Z			$\overline{W}$	
C_L	C_R	Fermions	C_L	C_R	Fermions
$\bar{g}^{\dagger}/2$	0	$\nu_e\,\bar{\nu}_e,\ \nu_\mu\,\bar{\nu}_\mu,\ \nu_\tau\,\bar{\nu}_\tau$	$\frac{g}{\sqrt{2}}$	0	$\bar{e}\,\bar{\nu},\bar{\mu}\,\bar{\nu},\bar{\tau}\,\bar{\nu}$
$(-\frac{1}{2} + \overset{**}{X_w})\,\bar{g}$	$X_w\,\bar{g}$	$e^+\,e^-,\ \mu^+\,\mu^-,\ \tau^+\,\tau^-$	$\frac{g}{\sqrt{2}}\cos\overset{*}{\theta_c}$	0	$\bar{u}\,d,\ \bar{c}\,s$
$(-\frac{1}{2} - \frac{2}{3}\,X_w)\,\bar{g}$	$-\frac{2}{3}\,X_w\bar{g}$	$u\,\bar{u},\ c\,\bar{c},\ t\,\bar{t}$	$\frac{g}{\sqrt{2}}\sin\theta_c$	0	$\bar{u}\,s,\ \bar{c}\,d$
$(-\frac{1}{2} + \frac{1}{3}\,X_w)\,\bar{g}$	$\frac{1}{3}\,X_w\,\bar{g}$	$d\,\bar{d},\ s\,\bar{s},\ b\,\bar{b}$	$\frac{g}{\sqrt{2}}$	0	$b\,\bar{t}$ ††

* θ_c — kabbibo angle,
** $X_w = \sin^2\theta_w \simeq \frac{1}{4}$. † $\bar{g}^2/\,M_z^2 = g^2/\,M_w^2 = 4\,\sqrt{2}\,G_F$.
†† As there is no sufficient information about t and b quarks, we assume that their interaction with W-bosons is the same as those of the e and ν.



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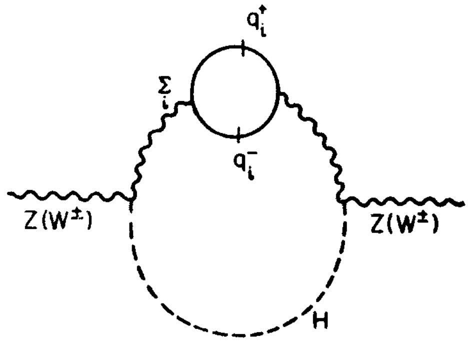

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Figure 2. Graph showing the contribution to I_m II coming from the process $Z(W^{\pm})\rightarrow$ all possible fermion pairs (quarks and leptons). $q_i^{\pm}$ denote the four-momenta of fermion pairs in the final state and the summation runs over all possible decay channels of $Z(W^{\pm})$ boson in the Weinberg-Salam theory. Dashes on the propagators denote that they are taken on the mass surface.

Figure 3. Graph showing the contribution to I_m II coming from the process $Z(W^{\pm})\rightarrow H_+$ all possible fermion pairs. Other notations are the same as in figure 2.

$$F_2(\alpha) = \frac{K^2/M^2 \cdot \sum_i [(C_L^i)^2 + (C_R^i)^2]}{(2\pi)^3\,(24)^2}$$

$$\times \int_0^{(\sqrt{\alpha}-\sqrt{\gamma})^2} \frac{[8\alpha\beta + (\alpha+\beta-\gamma)^2]\,\sqrt{(\alpha+\beta-\gamma)^2 - 4\alpha\beta}}{\alpha^2\,[(\beta-1)^2 + (F_1(\beta))^2]}\,\mathrm{d}\beta,$$

where

$$\alpha = Q^2/M^2,\ \beta = (q^+ + q^-)/M^2,\ \gamma = (M_H/M)^2,$$

$$K = 2\,(\sqrt{2}\,G_F)^{1/2}.\ M^2.$$

The total section σ_T will look like

$$\sigma_T(\alpha) = \frac{[(C_L^j)^2 + (C_R^j)^2]\,[F_1(\alpha) + F_2(\alpha)]}{D/2\,[(1-\alpha)^2 + (F_1(\alpha) + F_2(\alpha))^2]\cdot M^2} = \sigma_T(f_j\bar{f}_j\rightarrow), \tag{7}$$

$$\sigma_T(f_j\bar{f}_j\rightarrow) = \sigma_T^1 + \sigma_T^2 = \sigma_T^1\,(f_j\bar{f}_j\rightarrow Z\,(W^{\pm})\rightarrow \text{all fermions})$$

$$+\ \sigma_T^2\,(f_j\bar{f}_j\rightarrow Z\,(W^{\pm})\rightarrow H + \text{all fermions}).$$

The factor D in (7) appears as a result of averaging over spin and colour states of the fermions in the initial state. In particular, if $f_j\bar{f}_j = e^+\,e^-$ we at once recover the formula for $e^+ e^-$ collisions with $D = 4$, $C_L^2 + C_R^2 = \bar{g}^2/8$. (Sami and Fainberg 1981a). In the same way we can get from (7) the formula for $e\nu$ collisions (Sami and Fainberg 1982).

From (7), the expression for the total cross-section of process (3) is

$$\sigma_T^2 = \frac{F_2(\alpha)}{D/2[(1-\alpha)^2 + (F_1(\alpha) + F_2(\alpha))^2] \cdot M^2} \tag{8}$$

In the zero width approximation for $Z(W)$ boson ($F_1 \to 0$), its propagator (in our notations) can be replaced by $\pi \cdot \delta(\beta - 1)/F_1(1)$. The latter immediately reduces δ_T^2 to

$$\sigma(f_i \bar{f}_j \to H Z(W^{\pm})),$$

i.e. the cross-section of (3) calculated assuming the $Z(W^{\pm})$ boson to be stable. Had we considered some appropriate $Z(W^{\pm})$ boson decay channel, the effect of adding on the decay would have been the multiplication of the total cross-section by the branching ratio for $Z(W^{\pm})$ into that very channel.

It should be noted that process $f_i \bar{f}_j \to HZ\ (W^{\pm})$ is allowed for $\sqrt{\alpha} \geqslant 1 + \sqrt{\gamma}$, however, the process (3) goes for $\alpha \geqslant \gamma$. At $\alpha = 1$ ($M_H < M$, $r \leqslant \alpha < (1+\sqrt{r})^2$), σ_T^2 takes the maximum value. If $M_H \geqslant M$ the cross-section is very small until $\sqrt{\alpha} \geqslant 2$. In this case, most of the contribution to the integral over β in (8) comes from the neighbourhood of the point $\beta = 1$ ($F_1(1) \simeq 2{\cdot}10^{-2}$) which now lies in the domain of integration. In the neighbourhood of $\beta = 1$, we have the approximation:

$$\int \frac{\psi(\beta)}{[(1-\beta)^2 + (F_1(\beta))^2]}\, d\beta \simeq \frac{\psi(1)\,\pi}{F_1(1)}, \tag{9}$$

where $\psi(\beta)$ is some function continuous in the neighbourhood of $\beta = 1$. Approximation (9) is clearly equivalent to replacing the propagator by the delta-function. Thus in this case the zero width approximation can be used. It turns out that for $\sqrt{\gamma} \leqslant \sqrt{\alpha} < (1 + \sqrt{\gamma})$, σ_T^2 cannot be reduced to $\sigma(f_i \bar{f}_j \to H Z(W^{\pm}))$.

The above analysis also holds true for hadron-hadron collisions, as in this case, the cross-sections are expressed through σ_T^2.

3. Hadronic production of Higgs bosons.

In § 2, we have done the background calculations for the processes (1) and (2). The elementary cross-section σ_T taken with the structure functions and being integrated over the appropriate variables yields the desired result for *pp* and $p\bar{p}$ collisions. In fact, the total cross-section of all the processes, going through $Z(W)$ boson in *pp* and $p\bar{p}$ colliding beams, in the parton model is given by the expression (Quigg 1977; Peierls *et al* 1977; Brown and Mikaelian 1979)

$$\sigma_T(hh' \to) = \sum_{\substack{j \\ \text{flavours}}} \int \int_{\text{threshold}} dx_h\, dx_{h'}\, [q_j^h(X_h)\bar{q}_j^{h'}(X_{h'})$$

$$+ \bar{q}_j^h(X_h)\, q_j^{h'}(X_{h'})] \times \sigma_T(f_j \bar{f}_j \to), \tag{10}$$

$$\sigma_T(hh' \to) = \sigma_T(hh' \to l^+ l^- X) + \sigma_T(hh' \to l^+ l^- X H),$$

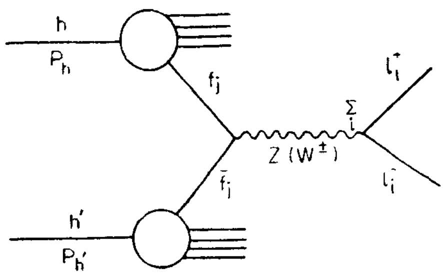
where $l^+ l^- = \sum_i l_i^+ l_{i-}^+$ summation is over all the decay channels into leptons, $hh' = pp\,(p\bar{p})$, $q_j^h(X_h)$ is the probability of finding the jth quark in hadron h with the momentum fraction X_h. $\sigma_T\,(f_j \bar{f}_j \to)$ has been prepared in § 1. Here we have used the Drell-Yan approximation (Drell and Yan 1970) for the production of $Z(W^\pm)$ bosons in pp, $p\bar{p}$ collisions (figures 4 and 5). The total cross-section of the process that we are interested in (figure 5), can be written as

$$\sigma_T(hh' \to H X l^+ l^-) = \frac{A_{Z,W} \cdot K^2/M^2}{4M^2 \cdot (2\pi)^3 (24)^2} \cdot \int_\gamma^\delta \int_{\theta/\delta}^1 \Phi_{z,w}(X_h, \theta) f(\theta)\, d\theta\, dX_h. \tag{11}$$

where $M^2\,\delta = (P_h + P_{h'})^2$ c.m. squared energy of hh',

$\theta = \delta \cdot X_h \cdot X_{h'}$, $A_Z = \frac{3}{2}\, g^4$,

$A_W = \frac{9}{8}\, g^4$.



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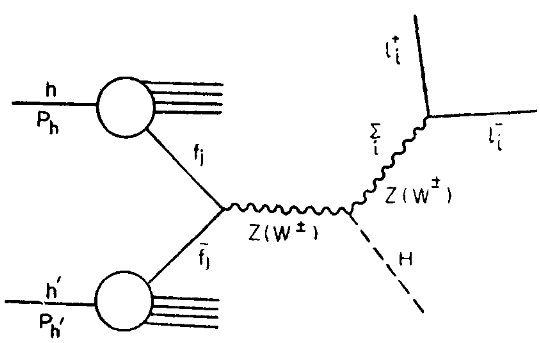

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Figure 4. Drell-Yan mechanism for lepton-pair production. P_h, $P_{h'}$ denote the four-momenta of hadrons h and h' respectively. $f_j\,\bar{f}_j$ is the appropriate quark (fermion) pair annihilating into $Z(W^\pm)$ boson. Summation runs over all possible lepton-pairs into which $Z(W^\pm)$ boson may decay.

Figure 5. Drell-Yan mechanism for lepton-pair and Higgs boson production. Other notations are the same as in figure 4.

The structure functions for PP collisions look like (we use the approximation $\theta_c = 0$)

$$\Phi_Z = \frac{2}{3\delta \cdot X_h} \Big\{ [u^-(X_h)\, u\,(\theta/\delta X_h) + (X_h \longleftrightarrow \theta/\delta X_h)]$$

$$\times \left(\frac{1}{4} - \frac{2X_W}{3} + \frac{8}{9} X_W^2\right) + [\bar{d}\,(X_h)\, d\,(\theta/\delta X_h) + (X_h \longleftrightarrow \theta/\delta X_h)$$

$$+ \bar{S}\,(X_h)\, S\,(\theta/\delta X_h) + (X_h \longleftrightarrow \theta/\delta X_h)] \left(\frac{1}{4} - \frac{X_W}{3} + \frac{2}{3} X_W^2\right)\Big\},$$

$$\Phi_{W^+} = \frac{1}{3\,\delta\, X_h}\, [u\,(X_h)\, \bar{d}\,(\theta/\delta \cdot X_h) + (X_h \longleftrightarrow \theta/\delta\; X_h)],$$

$$\Phi_{\overline{W}} = \frac{1}{3\,\delta\, X_h}\, [u^-(X_h)\, d\,(\theta/\delta\; X_h) + (X_h \longleftrightarrow \theta/\delta\; X_h)],$$

$$f(\theta) = \int_0^{(\sqrt{\theta} - \sqrt{\gamma})^2} \frac{[8\theta\beta + (\theta + \beta - \gamma)^2]\, \sqrt{(\theta + \beta - \gamma)^2 - 4\theta\beta}}{\theta^2\, [(\theta - 1)^2 + (F_1\,(\theta) + F_2\,(\theta))^2]\; [(\beta - 1)^2 + (F_1\,(\beta))^2} \, d\beta.$$

The parton distributions in proton are known (Okada *et al* 1976). For antiproton, read $\bar{u} \longleftrightarrow u$, $\bar{d} \longleftrightarrow d$. As seen from (11), the distribution $d\sigma/d\theta\; (hh' \to H\, X\, l^+ l^-)$ has a sharp peak about $\theta = 1$. The presence of such a resonance in the distribution is because of, an unstable particle in the intermediate state. In view of the fact, that Glashow *et al* have neglected the $Z(W^{\pm})$ boson decay effects, their distribution over θ has no peaks. The peak about $\theta = 1$ in the distribution is always attained for $\gamma < 1$ ($M_H < M$). Therefore, the integration over θ $r \leqslant \theta < (1 + \sqrt{r})$ in (11) can be evaluated with the help of narrow width approximation (9). The integral left after the approximation (9) is calculated numerically. If $M_H > M$, there is little hope to detect H-boson in the nearest future. In such a case the $Z(W)$ boson will be detected before the discovery of H-boson. This is due to the weak coupling of Higgs particles with matter (leptons and quarks). The association of $Z(W^{\pm})$ with H also provides a weak signal for H-boson production, *e.g.*

$$\sigma\,(pp \to H\; X\; Z) \simeq 10^{-39}\text{cm}^2,$$

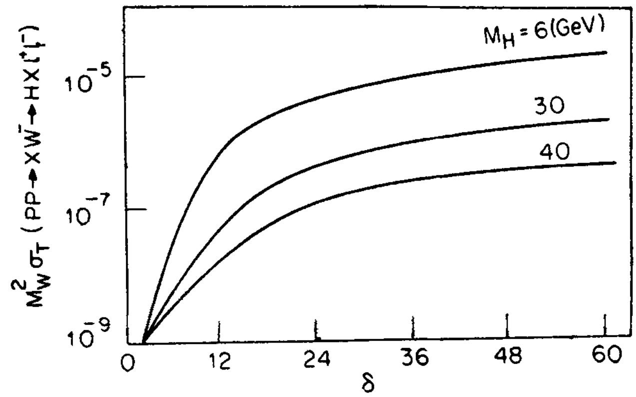
for $\sqrt{S} \simeq 1000$ GeV; $M_H \simeq 150$ GeV.

Note that for $\theta \geqslant (1 + \sqrt{r})^2$ our cross-sections are reduced to those of Glashow *et al* multiplied by the branching ratio for $Z\,(W^{\pm})$ boson into l^+l^- channel and this case has been omitted here.

The total cross-sections for pp, $p\bar{p}$ are plotted in figures 6, 7, 8, 9 and 10 for different H-boson masses. As seen from the figures all reactions start from $\delta = 1$. In fact, the collisions take place even for $\gamma \leqslant \delta < 1$ but the cross-sections are very small at such low energies. It is also seen from the figures that the cross-sections are big enough for $\delta\; (S/M^2) \simeq 60$. So, if the H-boson is not too heavy ($M_H < M$), one can hope to detect it in the near future. The numerical values of calculated cross-sections are of the same order as of those estimated by Glashow *et al e.g.*,

$$\frac{\sigma\,(p\bar{p} \to ZH + X)}{\sigma\,(p\bar{p} \to l^{+}l^{-}H + X)} \simeq 4,$$

for $\sqrt{S} = 700$ GeV and $M_H = 30$ GeV.



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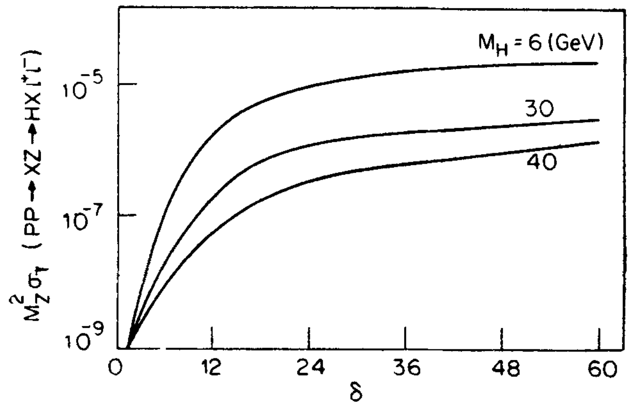

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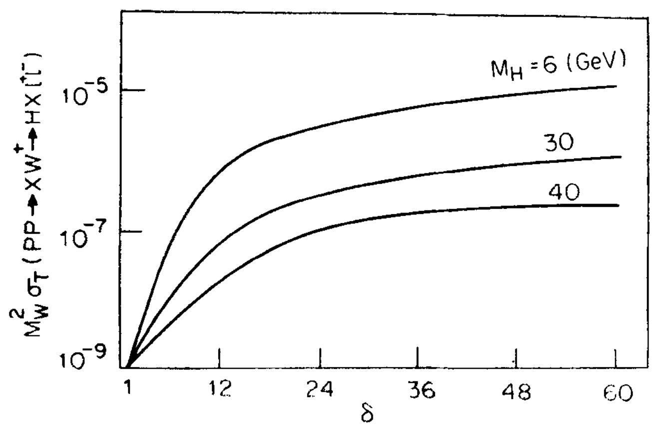

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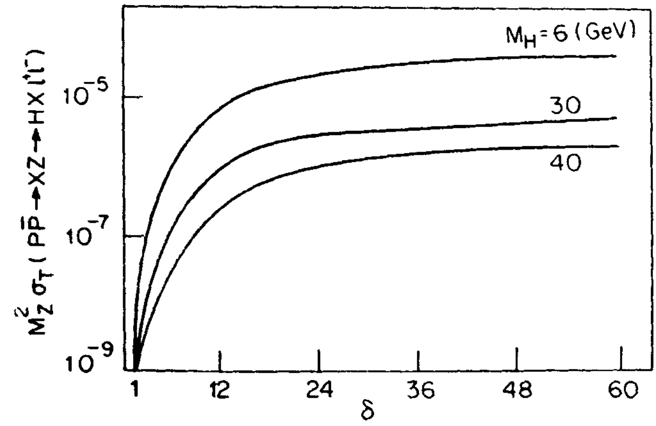

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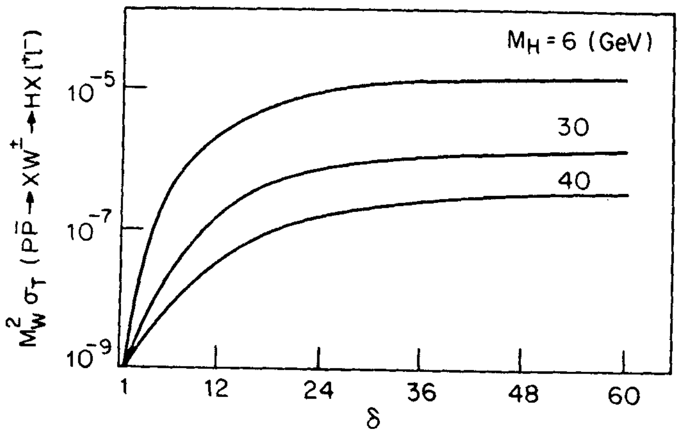

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Figures 6–10. Figures 6–10 plot of $M^2\,\sigma_T$ vs. δ (S/M^2) for three different values of Higgs boson mass. S is the total energy squared of pp $(p\bar{p})$ in the centre-of-mass system. $[\sigma_T] = [1/M^2]$.

If $M_H < M$, $\sigma\,(hh' \to HX\,l^+l^-)$ cannot be reduced to $\sigma\,(hh' \to HXZ\,(W^{\pm}))$. In this case, one should consider $Z(W^{\pm})$ decay and calculate $\sigma\,(hh' \to HXl^+l^-)$ instead of $\sigma\,(hh' \to HXZ\,(W^{\pm})$.

Finally, it should be noted that for better accuracy, one should go beyond the simple parton model and use the more precise theory that is to say QCD. In the leading order, in QCD all parton model formulas are left unchanged with an exception that the structure functions depend on Q^2 and the latter is calculable. The importance of higher order corrections has also been emphasized in the literature. This theory has not only strong theoretical background but also provides better experimental predictions (Buras 1980). The QCD corrections to the parton model may also seem to be important for the problem discussed in this paper.

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Notations used in the paper

ν_e —Electronic Neutrino
ν_μ —Muonic Neutrino
ν_τ —Taonic Neutrino
τ^- —Tao Lepton
τ^+ —Antitao Lepton
μ^+ —Antimuon
μ^- —Muon
$\Pi(Q^2)$ —*Pi* (Polarisation operator)