Preprint ЕФИ-850(I)-86

ԵՐԵՎԱՆԻ ՖԻԶԻԿԱՅԻ ԻՆՍՏԻՏՈՒՏ
ЕРЕВАНСКИЙ ФИЗИЧЕСКИЙ ИНСТИТУТ

R.Sh. EGORIAN

STANDARD HIGGS BOSON DISTRIBUTION OVER $q_\perp$ BORN IN HADRON-HADRON COLLISIONS

ЦНИИатоминформ

ЕРЕВАН-1986

Նախնատիպ ЕФИ–850(I)–86

Ռ.Շ.ԵԳՈՐՅԱՆ

ՀԱԴՐՈՆ-ՀԱԴՐՈՆԱՅԻՆ ԲԱԽՈՒՄՆԵՐՈՒՄ ԾՆՎԱԾ ՍՏԱՆԴԱՐՏ ՀԻԳՍԻ ԲՈԶՈՆԻ ԲԱՇԽՈՒՄԸ ԸՍՏ

Ուսումնասիրվում է ստանդարտ Հիգսի H_0 բոզոնի լայնական $q_\perp$ իմպուլսով սպեկտրը՝ կազմավորման տարբեր կանալներում. $q\bar{q}$-կանալ, $gq(\bar{q})$-կանալ, gg-կանալ, հադրոնային բախումներում: Ցույց է տրված, որ առկա են որակական տարբերություններ $d\sigma/dq_\perp^2$ H_0-բոզոնի և $\chi^{(0)}_{QQ}$-$(Q=b,t\ldots)$-մեզոնի միջև, օրինակ, քվարկային կանալում: Դա թույլ է տալիս միարժեքորեն նույնացնել H_0-ն $d\sigma/dq_\perp^2$ սպեկտրով $M_H \geqslant 10$ ԳէՎ զանգվածների ամբողջ տիրույթում:

Քննարկվում է բարձր կարգի խոտորումների դերը H_0 բոզոնների $d\sigma/dq_\perp^2$ սպեկտրներում:

Երևանի ֆիզիկայի ինստիտուտ

Երևան 1986

1. Introduction

After $W^{\pm}$ and Z_0 bosons were discovered [1-4] the experimental search of Higgs scalar bosons, predicted by the model of Glashow-Weinberg- Salam (GWS) electroweak interaction [5-7] , became one of the top-priority problems of modern elementary--particle physics.

Despite the fact that the detection and identification of scalar neutral Higgs H_0 bosons of GWS minimal scheme is rather a complicated problem (due to their highly insignificant bond with ordinary quarks and leptons and also to their specific decay, leading to many-particle final states) it is presumed to continue the investigation of Higgs spectrum on now developing accelerators of superhigh energies of the next generation with e^+e^- and $pp(p\bar{p})$ colliding beams [8].

The velocity spectrum of standard H_0 boson produced at pp collisions had been studied in Ref. [9] where it was suggested that standard Higgs bosons might be produced due to quark-antiquark or gluon pair annihilation (H_0 boson-gluon interaction occurring at one-loop level [10], see also [9] and [11]). It was mentioned there that the gluonic mechanism essentially contributed into the cross-section of direct $pp(p\bar{p})$ production of H_0 boson (the contribution of the quark channel being minute due to small mass of u , d , s current quarks).

In Ref.[8] it was mentioned that the gluonic production of Higgs bosons with mass $M_H \geqslant 180$GeV is quite perspective in colliding pp or $p\bar{p}$ beams.

The characteristics of spectra over the transverse moment of standard Higgs H_o bosons born at hadron-hadron collisions, are discussed below. It will be shown that all three channels of production ($q\bar{q} \to H_o + g$), ($gq(\bar{q}) \to H_o + q(\bar{q})$) and ($gg \to H_o + g$) are essential for $d\sigma/dq_\perp^2$ spectra of H_o bosons. The $d\sigma/dq_\perp^2$ spectra of H_o bosons in different channels of production ($q\bar{q}$, $gq(\bar{q})$ and gg) at pp and $p\bar{p}$ collisions will be calculated.

2. The Influence of the Highest Corrections of Perturbation Theory on the Spectra of Heavy Higgs Bosons (Gluon form factor)

Unlike the masses of fermions and gauge bosons, the mass of H_o boson is not fixed and is the only free parameter in GWS model of electroweak interaction. In this model M_H is uniquely bound with Higgs-field interaction constant λ : $(\lambda/16) H_o^4 = (M_H^2/8)(G_F\sqrt{2}) H_o^4$ (G_F is the Fermi constant). This permits to obtain theoretical restrictions imposed on H_o boson mass in terms of possible values of λ . Such limitations to M_H , coming from requirements of stability of physical vacuum at $|\varphi| = M_H/(\lambda\sqrt{2})$ (φ is the Higgs doublet field) and from the possibility of perturbation theory application to Higgs boson, are as follows: 3.5GeV $\lesssim M_H \lesssim$ 1TeV [12-14] (see also Refs [8] and [11]). There are also other theoretical limitations on M_H from below: $M_H \gtrsim$ 260MeV – for metastable vacuum, $M_H \gtrsim$ 9.2GeV – with regard to cosmological aspects of Universe evolution [14] (see also [8]) Though, in future, we shall consider Higgs bosons with $M_H \gtrsim 10$GeV ($\alpha_s(M_H^2)/\pi \ll 1$). Let us name such Higgs bosons as "heavy" ones.

From the QCD point of view the heavy Higgs bosons hadroproduction is a hard process. Here the application of perturbation theory is guaranteed by small conversion time τ of protons into heavy H_0 bosons. But, as mentioned in paragraphs 1 and 3, the H_0 boson production due to quark-antiquark annihilation is overwhelmed by that due to gluon fusion $gg \to H_0$ (because of small mass of current quarks). And the latter process (see Fig.1c,d) being analogous to that of C-even states of heavy quarkonium $Q\bar{Q}$ (i.e. of χ meson) due to the fusion of two gluons (see Fig.1c,d). In Ref. [15] it was shown that the full cross-section of χ -meson production process is mainly collected from a wide range of transverse momenta $q_\perp^2 \ll q^2 (= M_\chi^2)$ where $d\sigma/dq_\perp^2 \sim \sigma \frac{\alpha_s}{q_\perp^2}$. This results from the lowest order of perturbation theory for χ -meson production, when $q_\perp$ of χ -meson is compensated by one parton (e.g. by a quark, as shown in Fig. 1c or by a gluon, as shown in Fig.1d). The logarithmic nature of spectrum is due to parton "quasi-reality" $K(|K^2| \sim q_\perp^2 \ll q^2)$. It means that the real measure of hardness of process at registration of $q_\perp$ is $q_\perp^2$ and not the full mass q^2 [15,16].

Two parametrically different scales of characteristic distances $1/q \ll 1/q_\perp$ arise in this problem. In consequence, at taking the highest orders of perturbation theory into account, the differential cross-section of χ -meson production obtains a factor: effective, double-logarithmic gluon form factor raised to the second power [15]

$$\frac{d\sigma}{dq_\perp^2} \sim 6\ \frac{\alpha_s}{q_\perp^2} T_G^2(q_\perp^2, q^2), \qquad (1)$$

where $T_G^2(q_\perp^2 q^2)$ has the form

$$T_G^2(q_\perp^2, q^2) = \exp\left\{-C_v \int_{q_\perp^2}^{q^2} \frac{dK_\perp^2}{K_\perp^2} \frac{\alpha_s(K_\perp^2)}{\pi} \ln\left(\frac{q^2}{q_\perp^2}\right)\right\}. \qquad (2)$$

Consequently, the same is true for H_o boson production. This is quite useful, as a wide distribution over $q_\perp$ (resulting from $T_G^2(q_\perp^2, q^2$) may alleviate the background conditions for H_o boson identification in future experiments on H_o boson search at hadronic reactions.

It should be noted that the only difference now is that $q^2 = M_{H_o}^2$ enters Eq.(2) instead of $q^2 = M_\chi^2$ in the case of χ production.

The aforesaid reveals a unique possibility for unambiguous identification of H_o boson beyond the range of masses $\chi_{Q\bar{Q}}$ ($M_H \neq M_{\chi_{Q\bar{Q}}}$, $Q = b, t, \ldots$) over the enlarged spectrum $d\sigma/dq_\perp^2$ at H_o production in hadron-hadron collisions.

It should be also noted that there might exist a qualitative difference between $d\sigma/dq_\perp^2$ spectra of H_o bosons and χ-mesons born at hadron-hadron collisions, due to non-relativistic nature of χ-mesons. This would permit an unambiguous identification of H_o bosons(though both, H_o and χ_c are 0^+ particles) even in the range of masses $\chi_{Q\bar{Q}}$ (i.e. at $M_H = M_{\chi_{Q\bar{Q}}}$, $Q = b, t, \ldots$).

3. The Quark Channel of Heavy H_o Boson Production

The investigation of H_o boson production quark channel

$q\bar{q} \to H_0 + g$ is an interesting task in itself, for at any collision energy $\sqrt{S}$ it determines the difference of cross-sections $[d\sigma(a b \to H_0 + \dots) - d\sigma(\bar{a} b \to H_0 + \dots)]$, where a and b are colliding hadrons (for instance, a and b protons).

The $d\sigma/dq_{\gamma\perp}^2$ cross-section of $q\bar{q} \to H_0 + g$ process is determined by a set of Feynman diagrams shown in Fig.2. As the vertex $q\bar{q}H_0$ (in the diagrams 2a-h) is proportional to small current masses of quarks (it refers to both, valence quarks u, d, and sea quarks u, d, s, while the content of heavy quarks c, b, .. is very small in the sea) and the vertex $H_0 \to 2g$ (in the diagram 2i), roughly speaking, is proportional to the mass of two gluons, though it contains a small factor $\alpha_s(M_H^2)/\pi$ (see appendix 1), then the diagram 2i will dominate in the cross-section $d\sigma/dq_\perp^2$ of the process $q\bar{q} \to H_0 + g$.

The differential cross-section of H_0 production with fast y and with transverse moment $q_{\gamma\perp}$, corresponding to the diagram 2i, has the form

$$\frac{d\sigma^{(q\bar{q})}}{dq_{\gamma\perp}^2 dy} = \frac{\alpha_s}{4S} \cdot \frac{N^2-1}{2N} \iint \frac{dx_1}{x_1} \frac{dx_2}{x_2} D_1^q(x_1) D_2^{\bar{q}}(x_2) \delta_+(\ell^2) \times \tag{3}$$
$$\times \frac{1}{\Delta^4} \left\{ -\frac{1}{4} Sp\left(\hat{K}_1 \gamma_\nu \hat{K}_2 \gamma_\mu\right) \right\} M^{\mu\lambda}_{H_0 \to 2g}(\Delta, -\ell) M^{\nu\lambda}_{H_0 \to 2g}(-\Delta, \ell),$$

where $M^{\alpha\beta}_{H_0 \to 2g}(q_1, q_2)$ is the amplitude of H_0 boson decay into two gluons with $q_{\gamma 1}$ and q_2 momenta (α and β are the corresponding vector indices, see Fig.3). Substituting the explicit form of the amplitude $H_0 \to 2g$ in (app.1.13) into Eq.(3) and relying on the fact that $\ell_\lambda M^{\nu\lambda}_{H_0 \to 2g}(-\Delta, \ell) = \ell_\lambda M^{\mu\lambda}_{H_0 \to 2g}(\Delta, -\ell) = 0$, and also on relations $\Delta = K_1 + K_2 = q + \ell$; $K_1 = x_1 P_1$; $K_2 = x_2 P_2$; $q = y_1 P_1 + y_2 P_2 + q_\perp$; $P_1^2 \approx P_2^2 \approx 0$; $K_\perp = -q_\perp$; $K_1 q_\perp = K_2 q_\perp = 0$;

$\sqrt{S} = \sqrt{2P_1P_2}$; one obtains

$$\frac{d\sigma(12\to q\bar{q}\to H_0+q)}{dq_\perp^2\,dy} = \frac{\alpha_s^3 N_H^2}{18\pi^2 S N^2}\left(G_F\sqrt{2}\right)\int_\xi^{z_{max}} dz\,\frac{(1-z)}{z}\times$$
$$\times\left[D_1^q\left(\xi+\frac{q_\perp^2 z}{S\xi(1-z)}\right)D_2^{\bar{q}}\left(\frac{z}{\xi}\right)+D_1^q\left(\frac{\xi}{z}\right)D_2^{\bar{q}}\left(\xi+\frac{q_\perp^2 z}{S\xi(1-z)}\right)+\right.$$
$$\left.+D_1^{\bar{q}}\left(\xi+\frac{q_\perp^2 z}{S\xi(1-z)}\right)D_2^q\left(\frac{\xi}{z}\right)+D_1^{\bar{q}}\left(\frac{\xi}{z}\right)D_2^q\left(\xi+\frac{q_\perp^2 z}{S\xi(1-z)}\right)\right]. \tag{4}$$

where $\xi=\sqrt{(M_H^2+q_\perp^2)/S}$, $z_{max}=(1-\xi)\xi S/(q_\perp^2+(1-\xi)\xi S)$,
N_H is the number of aromas of quarks with $m_Q > M_H/2$,
N is the $SU(N)$ group dimension.

Then, at $q_\perp\to 0$, without concretization of the forms of distribution functions, we obtain the following form for the H_0 boson production spectrum

$$\frac{1}{\sigma}\cdot\frac{d\sigma^{(q\bar{q})}}{dq_\perp^2\,dy} = \int_\xi^{z_{max}} dz\,\frac{(1-z)}{z} \tag{5}$$

It is seen from (5) that at $q_\perp\to 0$ (M_H is fixed)

$$\frac{d\sigma^{(q\bar{q})}}{dq_\perp^2} \backsim const \tag{6}$$

Such behaviour of H_0 boson spectrum $d\sigma/dq_\perp^2$ at $q_\perp\to 0$ sharply differs from that of χ_0 -meson which formally has a pole at $q_\perp\to 0$: $d\sigma/dq_\perp^2 \backsim q_\perp^{-2}$ [15].

Such a qualitative difference in $d\sigma/dq_\perp^2$ spectra of H_0 boson and χ_0-meson at $q_\perp\to 0$,due to non-relativistic nature of χ_0-meson (both of them are 0^+ particles), will, apparently, unambiguously identify H_0 boson in the whole range of masses of χ_0-meson (i.e. $M_H = M_{\chi_{Q\bar{Q}}}$, $Q = b, t, \dots$).

4. Quark-Gluon Channel of Heavy H_0 Production

Dominating diagrams for the cross-section $d\sigma/dq_\perp^2$ of the process $gq(\bar{q}) \to H_0 + q(\bar{q})$ are the ones shown in Fig.4.

H_0 boson production differential cross-section $d\sigma/dy\,dq_\perp^2$ shown in diagram 4a has the form

$$\frac{d\sigma^{gq(\bar{q})}}{dq_\perp^2\,dy} = \frac{\alpha_s}{4S}\,\frac{1}{2N^2}\iint \frac{dx_1}{x_1}\,\frac{dx_2}{x_2} D_1^g(x_1) D_2^{q(\bar{q})}(x_2)\,\delta_+(\ell^2)\times \qquad (7)$$

$$\times \frac{1}{K^4}\left\{\frac{1}{4} Sp(\hat{K}_2\gamma_\nu\,\hat{\ell}\gamma_\mu)\right\} M^{\mu\sigma}_{H_0\to 2g}(K_1,K)\, M^{\nu\sigma'}_{H_0\to 2g}(K_1,K)\, g^\perp_{\sigma\sigma'}$$

where x_1 and x_2 are portions of hadron momenta taken away by a gluon or a quark, respectively (i.e. $K_1 = x_1 P_1$, $K_2 = x_2 P_2$); is the tensor perpendicular to collision plane ($g^\perp_{\sigma\tau} = g_{\sigma\tau} - (P_{1\sigma}P_{2\tau} + P_{1\tau}P_{2\sigma})/(P_1P_2)$); $M^{\sigma\mu}_{H_0\to 2g}$ and $M^{\sigma'\nu}_{H_0\to 2g}$ are the amplitudes of H_0 boson two-gluon decay. Substituting the explicit form of the amplitude of $H_0 \to 2g$ in (app.1.13) into Eq.(7) and relying on the relations $K_1^2 \approx K_2^2 \approx 0$ ($P_1^2 \approx P_2^2 \approx 0$); $K_1^\nu g^\perp_{\sigma\nu} = 0$; $(K_\perp)^2 = -q_\perp^2$, one obtains the following form for spectrum $d\sigma/dq_\perp^2$ of the process $12 \to gq(\bar{q}) \to H_0 + q(\bar{q})$

$$\frac{d\sigma(12\to gq(\bar{q})\to H_0+q(\bar{q}))}{dy\,dq_\perp^2} = \frac{2\alpha_s^3}{9\pi^2 S}\cdot\frac{N_H^2(G_F\sqrt{2})}{N^2(N^2-1)}\cdot\frac{1}{q_\perp^2}\int_\xi^{z_{max}} dz\,\frac{1+(1-z)^2}{z^2(1-z)}\times \qquad (8)$$

$$\times\left(q_\perp^2 + M_H^2(1-z)\right)\left[D_1^g\left(\xi + \frac{q_\perp^2 z}{S\xi(1-z)}\right) D_2^{q(\bar{q})}\left(\frac{\xi}{z}\right) + D_1^{q(\bar{q})}\left(\frac{\xi}{z}\right) D_2^g\left(\xi + \frac{q_\perp^2 z}{S\xi(1-z)}\right)\right]$$

And then, at $q_\perp \to 0$, without specifying the forms of distribution functions, the form of the spectrum $d\sigma/dq_\perp^2$ of H_0 boson production in this channel will be:

$$\frac{1}{6}\cdot\frac{d\sigma^{(qg)}}{dq_\perp^2\,dy}=\frac{1}{q_\perp^2}\int\limits_{\xi}^{z_{max}}dz\;\frac{1+(1-z)^2}{z^2(1-z)}\left(q_\perp^2+M_H^2(1-z)\right). \qquad (9)$$

It is interesting that this spectrum is like the one for χ_0-meson in the same channel of production: $d\sigma^{(qg)}_{\chi_c}/dq_\perp^2 \backsim q_\perp^{-2}$ (see Ref. 17). It is due to the fact that in this channel the relativistic nature of χ_0-meson can not manifest itself, as none of the gluons coupled with χ_c can be a soft one.

5. Gluon Channel of Heavy H_0 Production

The full set of Feynman diagrams for $d\sigma/dq_\perp^2dy$ cross-section of the process $gg\to H_0+g$ is presented in Fig.5. The form of differential cross-section $d\sigma/dq_\perp^2dy$,which is the sum of all these diagrams, is

$$\frac{d\sigma^{gg}}{dq_\perp^2\,dy}=\frac{\alpha_s N}{4\pi(N^2-1)}\int\!\!\int\frac{dx_1}{x_1}\,\frac{dx_2}{x_2}\,D^g(x_1)\,D^g(x_2)\times \alpha_s(Q^2)\sum_{i=1}^{16}\Phi_i^{(gg)} \qquad (10)$$

[illegible]

$$\Phi_3^{(gg)} = -\frac{1}{4\Delta^4} M^{\mu\lambda}_{H_0\to 2g}(\Delta,-\ell) M^{\nu\lambda}_{H_0\to 2g}(-\Delta,\ell) \times$$

$$\times \Gamma_{\sigma\mu\tau}(K_1,-\Delta,K_2)\Gamma_{\sigma'\tau'\nu}(-K_1,-K_2,\Delta) g^{\perp}_{\sigma\sigma'} g^{\perp}_{\tau\tau'} \; ; \qquad (10c)$$

$$\Phi_4^{(gg)} = -\frac{1}{4} M^{\sigma\lambda\tau}_{H_0\to 3g}(K_1,-\ell,K_2) M^{\sigma'\tau'\lambda}_{H_0\to 3g}(-K_1,-K_2,\ell) g^{\perp}_{\sigma\sigma'} g^{\perp}_{\tau\tau'} \; ; \qquad (10d)$$

$$\Phi_5^{(gg)} = -\frac{1}{4K^2K'^2} M^{\sigma\mu}_{H_0\to 2g}(K_1,K) M^{\tau'\nu}_{H_0\to 2g}(-K_2,-K') \times$$

$$(10e)$$

$$\times \Gamma_{\mu\lambda\tau}(-K,-\ell,K_2)\Gamma_{\nu\lambda\sigma'}(K',\ell,-K_1) g^{\perp}_{\sigma\sigma'} g^{\perp}_{\tau\tau'} \; ;$$

$$\Phi_6^{(gg)} = \Phi_5^{(gg)}(K_1 \leftrightarrow K_2,\; K \leftrightarrow K') = \Phi_5^{(gg)} \; ; \qquad (10f)$$

$$\Phi_7^{(gg)} = \Phi_8^{(gg)} = -\frac{1}{4\Delta^2} M^{\mu\lambda}_{H_0\to 2g}(\Delta,-\ell) M^{\sigma'\tau'\lambda}_{H_0\to 3g}(-K_1,-K_2,\ell) \times$$

$$(10g)$$

$$\times \Gamma_{\sigma\mu\tau}(K_1,-\Delta,K_2) g^{\perp}_{\sigma\sigma'} g^{\perp}_{\tau\tau'}$$

$$\Phi_9^{(gg)} = \Phi_{10}^{(gg)} = -\frac{1}{4\Delta^2K^2} M^{\mu\lambda}_{H_0\to 2g}(\Delta,-\ell) M^{\nu\sigma'}_{H_0\to 2g}(-K_1,-K) \times \qquad (10h)$$

$$\times \Gamma_{\sigma\mu\tau}(K_1,-\Delta,K_2)\Gamma_{\tau'\lambda\nu}(-K_2,\ell,K) g^{\perp}_{\sigma\sigma'} g^{\perp}_{\tau\tau'}$$

$$(10i)$$

$$\Phi_{11}^{(gg)} = \Phi_{12}^{(gg)} = \Phi_{10}^{(gg)}(K_1 \leftrightarrow K_2,\; K \leftrightarrow K')$$

$$\Phi_{13}^{(gg)} = \Phi_{14}^{(gg)} = -\frac{1}{4K^2} M^{\sigma\mu}_{H_0\to 2g}(K_1,K)\, M^{\sigma'\tau'\lambda}_{H_0\to 3g}(-K_1,-K_2,\ell)\times \tag{10j}$$

$$\times\, \Gamma_{\mu\lambda\tau}(-K,-\ell,K_2)\, q^{\perp}_{\sigma\sigma'}\, q^{\perp}_{\tau\tau'}$$

$$\Phi_{15}^{(gg)} = \Phi_{16}^{(gg)} = \Phi_{14}^{(gg)}(K_1 \leftrightarrow K_2,\ K \leftrightarrow K') \tag{10k}$$

In all relations (10) $\Gamma_{\alpha\beta\gamma}$ is a simple three-gluon vertex.

Substituting the expressions for amplitudes $H_0 \to 2g$ and $H_0 \to 3g$ in (app.2.4) into Eq.(10) and then substituting (10) into (9) and doing simple transformations with account of all relations indicated in paragraphs 3 and 4, one will obtain the following expression for $d\sigma/dq_\perp^2$ in the gluon channel of production:

$$\frac{d\sigma(12\to gg\to H_0+g)}{dq_\perp^2\, dy} = \frac{\alpha_s^3 N_H^2 N (G_F\sqrt{2})}{9\pi^2 (N^2-1) S}\cdot\frac{1}{q_\perp^2}\int\limits_{\xi}^{z_{max}} dz\, D^g\!\left(\frac{\xi}{z}\right) D^g\!\left(\xi + \frac{q_\perp^2 z}{S\xi(1-z)}\right)\times \tag{11}$$

$$\times\left\{\left[2\left(q_\perp^2 + M_H^2(1-z)\right)\frac{1+(1-z)z^2}{z^2(1-z)} + \frac{M_H^2}{z(1-z)}\right] + \right.$$

$$\left. + q_\perp^2\left[\frac{q_\perp^2\left(2+7z(1-z)\right) + M_H^2(1-z)\left(2+4z(1-z)\right)}{z(1-z)\left(q_\perp^2 + M_H^2(1-z)\right)}\right]\right\}$$

It should be noted that, as it is seen from (11), the H_0 boson spectrum $d\sigma/dq_\perp^2$ also has a formal pole at $q_\perp \to 0$: $d\sigma/dq_\perp^2 \sim q_\perp^{-2}$, the explanation of which is just the same as in the case of quark-gluon channel of H_0 boson production (see paragraph 4).

The author is thankful to Yu.L. Dokshitzer and Ed.Sh. Egoryan for stimulating discussions and permanent attention towards this work and also to A.A. Anselm, S.G. Grigoryan and S.G. Matinyan for useful discussions.

Appendix 1. The Amplitude of Two-Gluon Decay of Standard H_0 Boson

The most general form of the amplitude of standard Higgs boson two-gluon decay, satisfying the conditions of Lorentz-invariance and C-invariance, is

$$M^{\alpha\beta;\,ab}_{H_0\to 2g}(q_1,q_2) = \left[A g^{\alpha\beta} + B\left(q_1^\alpha q_2^\beta + q_1^\beta q_2^\alpha\right) + C\left(q_1^\alpha q_1^\beta + q_2^\alpha q_2^\beta\right)\right]\delta_{ab} \qquad \text{(app.1.1)}$$

where α, β are the vector indices of gluons with momenta q_1 and q_2, respectively; δ_{ab} shows that over colour indices a, b of two gluons the amplitude of two-gluon decay of H_0 boson is a singlet.

It follows from the law of conservation of current that

$$q_1^\alpha e_2^\beta M^{\alpha\beta}_{H_0\to 2g}(q_1,q_2) = \left\{A q_1^\beta + B\left[q_1^2 q_2^\beta + (q_1 q_2) q_1^\beta\right] + C\left[q_1^2 q_1^\beta + (q_1 q_2) q_2^\beta\right]\right\} e_2^\beta = 0 \qquad \text{(app.1.2)}$$

$$q_2^\beta e_1^\alpha M^{\alpha\beta}_{H_0\to 2g}(q_1,q_2) = \left\{A q_2^\alpha + B\left[q_2^2 q_1^\alpha + (q_1 q_2) q_2^\alpha\right] + C\left[q_2^2 q_2^\alpha + (q_1 q_2) q_1^\alpha\right]\right\} e_1^\alpha = 0 \qquad \text{(app.1.3)}$$

Taking $e_1^\alpha q_1^\alpha = e_2^\beta q_2^\beta = 0$ (gluons with physical polarization), one obtains

$$q_1^{\beta} e_2^{\beta} \left\{ A + B(q_1 q_2) + C q_1^2 \right\} = 0 \qquad \text{(app.1.4)}$$

$$q_2^{\alpha} e_1^{\alpha} \left\{ A + B(q_1 q_2) + C q_2^2 \right\} = 0 \qquad \text{(app.1.5)}$$

It follows from (app.1.4) and (app.1.5) that $C = 0$ and $A = -B/(q_1 q_2)$. Taking $B = -\lambda$ (to make λ dimensionless we devide it by M_H) we shall obtain the following form for the amplitude $H_0 \to 2g$:

$$M^{\alpha\beta;ab}_{H_0 \to 2g}(q_1, q_2) = \delta_{ab} \frac{\lambda}{M_H} \left(g^{\alpha\beta}(q_1 q_2) - q_1^{\alpha} q_2^{\beta} - q_1^{\beta} q_2^{\alpha} \right) \qquad \text{(app.1.6)}$$

Such representation of the amplitude of H_0 boson two-gluon decay, in fact, is equivalent to H_0 boson decay into gluons in the framework of local effective Lagrangian (see e.g., [11]). Therefore, using the width $\Gamma(H_0 \to 2g)$, calculated with the help of quark loops (see Fig.3), one may find the coefficient λ in (app.1.6) expressing λ in terms of H_0 boson two-gluon decay [11]

$$\Gamma(H_0 \to 2g) = N_H^2 \left(\frac{\alpha_s(M_H^2)}{\pi} \right)^2 \frac{G^2}{72\pi} M_H^3 \qquad \text{(app.1.7)}$$

where N_H is the number of different sorts of heavy quarks ($m_q > M_H/2$). Let us express λ in terms of $\Gamma(H_0 \to 2g)$, proceeding from relation $M_H \Gamma(H_0 \to 2g) = 2\,\mathrm{Im}\,\Sigma$ where (see Fig.6)

$$2\,\mathrm{Im}\,\Sigma = \frac{1}{2} \int M^{\alpha\beta;ab}_{H_0 \to 2g}(q_1, q_2) M^{\alpha\beta;ab}_{H_0 \to 2g}(q_1, q_2)^{*} \qquad \text{(app.1.8)}$$

$$\times (2\pi)\delta(q_1^2)(2\pi)\delta(q_2^2)\frac{d^4 q_1}{(2\pi)^4}$$

(The coefficient 1/2 prior to the integral comes from the identity of gluons).

Substituting $M^{\alpha\beta;ab}_{H_0\to 2g}(q_1, q_2)$ in (app.1.6) into (app.1.8) we have

$$2\,\mathrm{Im}\,\Sigma = \frac{1}{4(2\pi)^2}\lambda^2 M_H^2 (N^2-1)\int \delta(q_1^2)\,\delta(q^2-2q_1 q)\,d^4 q_1\,, \quad \text{(app.1.9)}$$

where N is the number of colours, and $q \equiv q_1 + q_2$.

Substituting $d^4 q_1 = |\vec{q}_1|^2 d|\vec{q}_1| dq_1^0\, d\Omega$ into (app.1.9) and passing into the system of centre of mass (i.e. $\vec{q}_1 = -\vec{q}_2$; $q_1^0 = q_2^0 = \frac{q^0}{2} = \frac{M_H}{2}$) one obtains

$$2M_H \cdot \Gamma(H_0 \to 2g) = 2\,\mathrm{Im}\,\Sigma = \frac{\lambda^2 M_H^2 (N^2-1)}{32\pi}\,. \quad \text{(app.1.10)}$$

Applying (app.1.7) to $\Gamma(H_0\to 2g)$ one obtains

$$N_H^2\left(\frac{\alpha_s}{\pi}\right)^2 \frac{v^2}{72\pi} M_H^3 = \frac{\lambda^2 M_H (N^2-1)}{64\pi}\,. \quad \text{(app.1.11)}$$

From (app.1.11) we get the following form for λ :

$$\lambda = \frac{2}{3} N_H \frac{\alpha_s}{\pi} v M_H \cdot \sqrt{\frac{2}{N^2-1}} \quad \text{(app.1.12)}$$

Substituting (app.1.12) into (app.1.6) one will obtain the following form for the amplitude of two-gluon decay of :

$$M^{\alpha\beta;ab}_{H_0\to 2g}(q_1,q_2) = \delta_{ab}\frac{2}{3}N_H v\sqrt{\frac{2}{N^2-1}}\left\{g^{\alpha\beta}(q_1 q_2) - q_1^\alpha q_2^\beta - q_1^\beta q_2^\alpha\right\}\frac{\alpha_s}{\pi} \quad \text{(app.1.13)}$$

Appendix 2. The Amplitude of Three-Gluon Decay of Standard H_0 Boson

The form of three-gluon vertex of H_0 boson decay can be obtained from the law of conservation of current for the sum of diagrams shown in Fig.7, using the form of vertex of H_0 two-gluon decay in (app.1.13).

The law of conservation of current for the sum of diagrams shown in Fig.7 has the form

$$e_1^{\alpha} e_2^{\beta} \ell^{\gamma} \Big\{ M^{\alpha\gamma\beta;\, acb}_{H_0\to 3g}(q_1, \ell, q_2) +$$

$$+ M^{\beta\delta,\, bd}_{H_0\to 2g}(q_2, q_1+\ell)\, i f_{acd}\, \Gamma_{\alpha\gamma\delta}(q_1, \ell, -q_1-\ell)\frac{1}{(q_1+\ell)^2} + \qquad \text{(app.2.1)}$$

$$+ M^{\alpha\delta,\, ad}_{H_0\to 2g}(q_1, q_2+\ell)\, i f_{cbd}\, \Gamma_{\gamma\beta\delta}(\ell, q_2, -q_2-\ell)\frac{1}{(q_2+\ell)^2} +$$

$$+ M^{\gamma\delta,\, cd}_{H_0\to 2g}(\ell, q_1+q_2)\, i f_{abd}\, \Gamma_{\alpha\beta\delta}(q_1, q_2, -q_1-q_2)\frac{1}{(q_1+q_2)^2} \Big\} = 0$$

Using the explicit form of the vertex of H_0 boson two-gluon decay in (app.1.13) and bearing in mind that the two-gluon vertex $H_0 \to 2g$ is a singlet over colour (i.e. substitute d in (app.2.1) by b, a, c in the second, the third and the fourth terms, correspondingly) and also substituting the familiar expressions of standard three-gluon vertices $\Gamma_{\alpha\gamma\delta}$, $\Gamma_{\gamma\beta\delta}$, $\Gamma_{\alpha\beta\delta}$ one will obtain the form of the amplitude of H_0 boson three-gluon decay

$$M^{\alpha\gamma\beta;ac\beta}_{H_0\to 3g}(q_1,\ell,q_2)=if_{ac\beta}\frac{\lambda}{M_H}\Gamma_{\alpha\gamma\beta}(q_1,\ell,q_2) \qquad \text{(app.2.2)}$$

where $\Gamma_{\alpha\gamma\beta}$ is a standard three-gluon vertex having the form

$$\Gamma_{\alpha\gamma\beta}(q_1,\ell,q_2)=\left\{g_{\alpha\gamma}(\ell-q_1)_\beta+g_{\gamma\beta}(q_2-\ell)_\alpha+g_{\beta\alpha}(q_1-q_2)_\gamma\right\} \qquad \text{(app.2.3)}$$

It should be noted that (app.2.2) has been obtained with regard to the fact that gluons are "real" ($e_1^\alpha q_{1\alpha}=e_2^\beta q_{2\beta}=0$)and have no mass ($q_1^2=q_2^2$ +=0). Substituting the expression for λ in (app.1.12) into (app.2.2) one obtains

$$M^{\alpha\gamma\beta;ac\beta}_{H_0\to 3g}(q_1,\ell,q_2)=if_{ac\beta}\,\frac{2}{3}N_H\left(\frac{\alpha_s}{\pi}\right)v\sqrt{\frac{2}{N^2-1}}\,\times$$

$$\times\left\{g_{\alpha\gamma}(\ell-q_1)_\beta+g_{\gamma\beta}(q_2-\ell)_\alpha+g_{\beta\alpha}(q_1-q_2)_\gamma\right\} \qquad \text{(app.2.4)}$$

Here (in the text too) $v=(G_F\sqrt{2})^{1/2}\approx(246)^{-1}$; f_{ijk} are the structure constants of the group SU(N) .

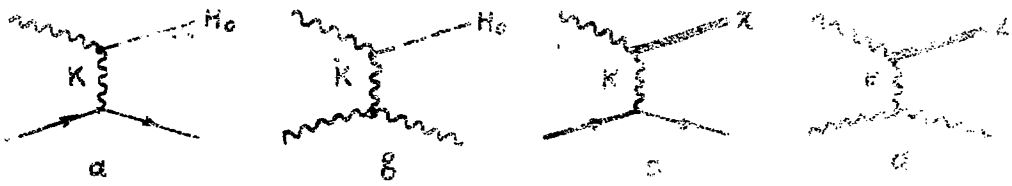
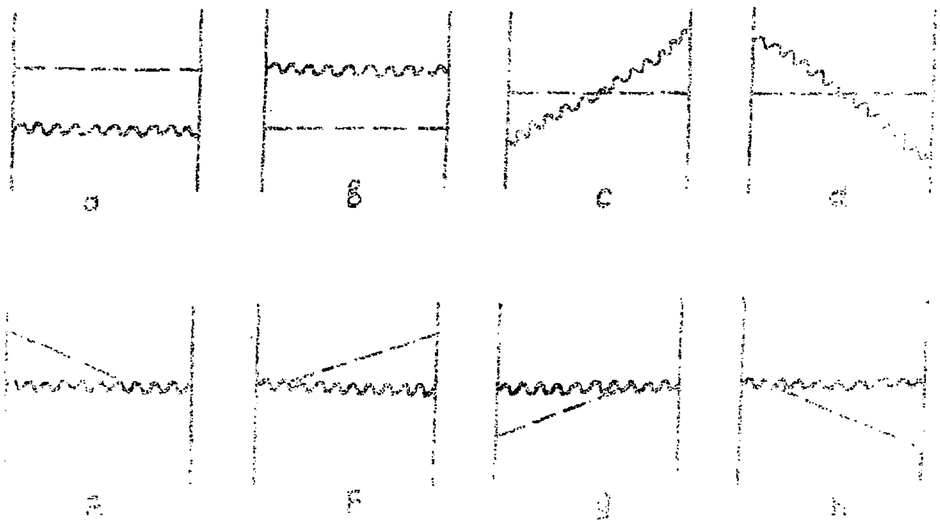
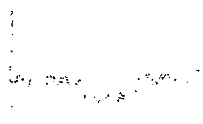


Fig. 1 The lowest order diagrams of H_0 production in qg-(a and gg-(б) channels; for the production of [illegible] the same channels.





[illegible] diagrams [illegible] production of [illegible]

[illegible] $q\bar{q} \to H_0$ [illegible]

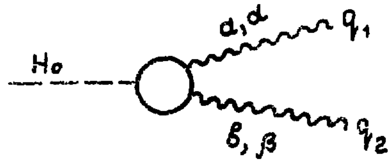


Fig.3 A diagram for the amplitude of H_o boson two-gluon decay.

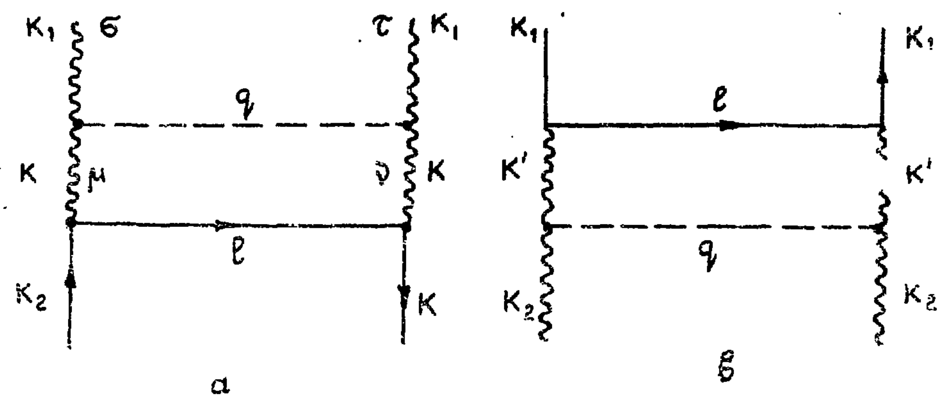


Fig.4 Dominating diagrams for cross-section $d\sigma/dq_\perp^2$ of process $gq(\bar{q}) \rightarrow H_o + q(\bar{q})$.

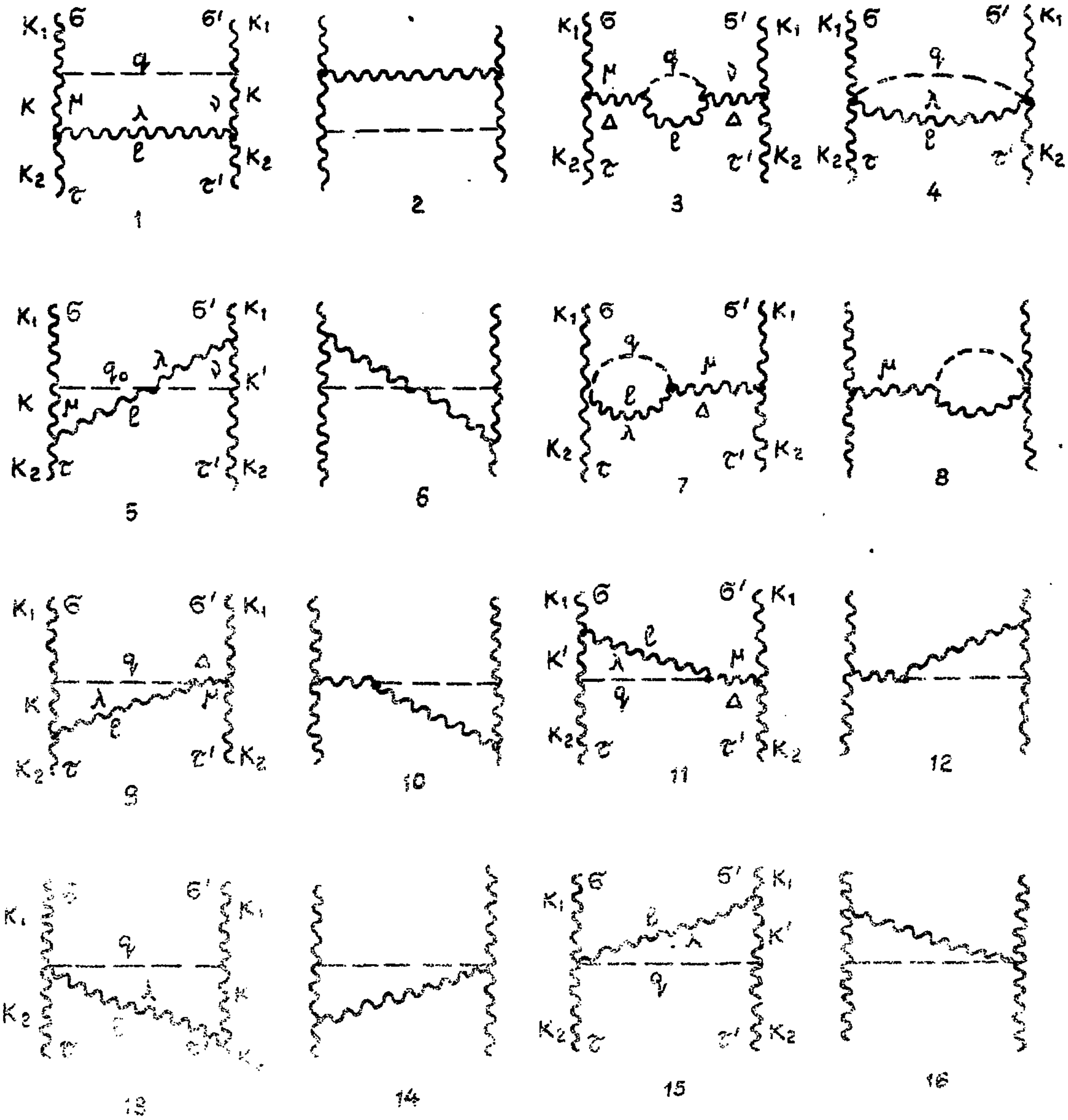


Fig.5 Full set of Feynman diagrams for cross-section $d\sigma/dq_{\perp}^{2}$ of process $gg \to H_0 + g$.

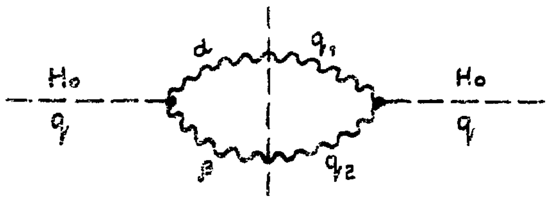


Fig.6 A diagram for determination of the width $\Gamma(H_o \to 2g)$.

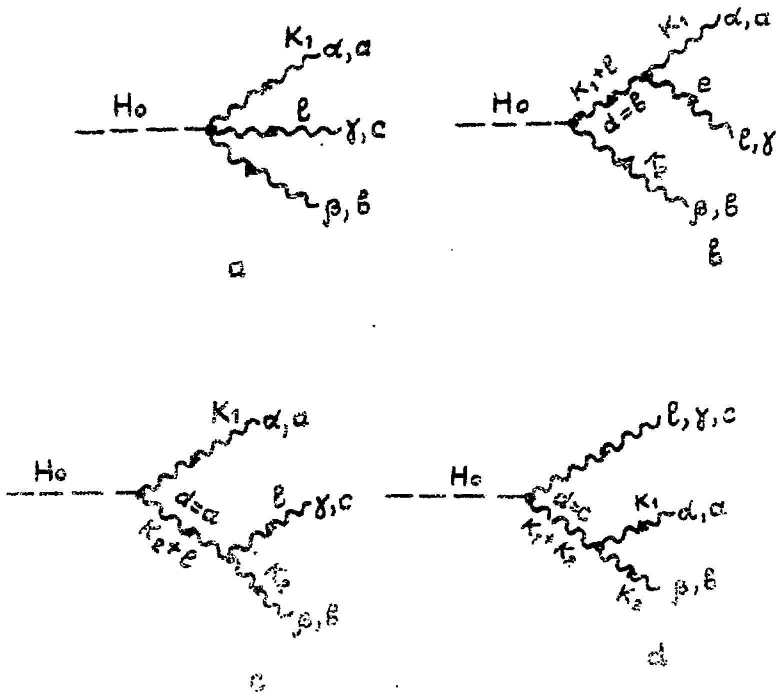


Fig.7 Diagrams providing the conservation of current.

REFERENCES

1 Arnison G. et al. (UA collab), Experimental observation of isolated large transverse energy electron with associated missing energy at s=540GeV, Phys.Lett., 1983, vol.122B, p.103;

Arnison G. et al. (UA collab), Further evidence for charged intermidiated vector bosons at the SPS collider, Phys.Lett., 1983, vol.129B, p.273.

2 Banner M., et al., (UA 2 collab), Observation of single isolated electrons of high transverse momentum in events widths missing transverse energy at the CERN pp collider, Phys.Lett., 1983, vol.122B, p.476.

3 Arnison G., et al. (UA1 collab), Experimental observation of lepton pairs of invariant mass around 95GeV/c^2 at the CERN SPS collider, Phys.Lett., 1983, vol.126B, p.398.

4 Bagnaia P., et al. (UA2 collab), Evidence for Zo → e^+e^- at the CERN pp collider, Phys.Lett., 1983, vol.129B, p.130.

5 Glashow S.L., Partial-Symmetries of weak interactions, Nucl. Phys., 1961, vol.22, p.579.

6 Weinberg S., A model of leptons, Phys.Rev.Lett., 1967, vol.19, p.1264.

7 Salam A., Proc. 8th Nobel Symposium, Stockholm 1968, ed.N. Svartholm (Almqvist and Wiksells, Stockholm, 1968).

8 Ансельм А.А., Уральцев Н.Г., Хозе В.А. Хиггсовские бозоны. Физика высоких энергий. Материалы XIX Зимней школы ЛИЯФ, 1984.

9 Georgi H.M., Glashow S.L., Hachaceh H.E., and NanopoulosD.V. Higgs bosons from two-gluon annihilation in Parton-Parton collisions, Phys.Rev.Lett., 1978, vol.40, p.692.
10 Wilczek F. Decays of heavy vector mesons into Higgs particles, Phys.Rev.Lett., 1977, vol.39, p.1304.
11 Окунь Л.Б. Лептоны и кварки. М.: Наука, 1981.
12 Линде А.Д. Динамическое восстановление симметрии и ограничений на массы и константы связи в модели Хиггса.Письма в ЖЭТФ, 1976, т.23, с.73.
13 Weinberg S., Mass of the Higgs Boson, Phys.Rev.Lett., 1976 vol.36, p.294.
14 Linde A.D. Vacuum instability, cosmology and constraints on particle mass in the Weinberg-Salam model, Phys.Lett., 1980, vol.92B, p.119.
15 Yu.L. Dokshitzer, R.Sh. Egorian, On peculiarities of formation of C-even heavy quarkonia in hadron-hadron collisions, prepr. EPI-692(7)-84; Докшицер Ю.Л., Егорян Р.Ш. Об особенностях образования С-четных тяжелых кваркониев в адронных соударениях. ЯФ, 1984, т.12, с.1547.
16. Yu.L. Dokshtzer, D.I. Dyahonov, S.I.Troyan, Hard Semi-inclusive processes in QCD, Phys.Lett., 1978, vol.78B, p.290.
17 R.Sh. Egorian, Peculiarities of the quark-gluon mechanism of the formation of C-even heavy quarkonia in hadron collisions, prepr. EPI-806(33)-85.

The manuscript was received 26 October 1985

Р.Ш.ЕГОРЯН

РАСПРЕДЕЛЕНИЕ ПО $q_\perp$ СТАНДАРТНОГО ХИГГСОВСКОГО БОЗОНА, РОЖДЕННОГО В АДРОН-АДРОННЫХ СОУДАРЕНИЯХ.

(на английском языке, перевод Г.А.Папяна)

Редактор Л.П.Мукаян
Технический редактор А.С.Абрамян

Подписано в печать 4/III-86г. ВФ-06546 Формат 60х84/16
Офсетная печать.Уч.изд.л.1,0 Тираж 299 экз.Ц.15 к.
Зак.тип.№ 158 Индекс 3624

Отпечатано в Ереванском физическом институте
Ереван 36, Маркаряна 2

индекс 3. 24

ЕРЕВАНСКИЙ ФИЗИЧЕСКИЙ ИНСТИТУТ