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R.A.ALANAKYAN,S.G.GRIGORYAN,S.G.MATINYAN

HIGGS BOSON PRODUCTION IN HADRONIC DECAYS OF C-EVEN QUARKONIA

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HIGGS BOSON PRODUCTION IN HADRONIC DECAYS OF C-EVEN QUARKONIA

The differential characteristics and total widths for H^0-boson production in hadronic decays of C-even heavy quarkonia (η , χ_0 , χ_1 , χ_2) are studied. A comparison of differential distributions and total widths for scalar and pseudoscalar boson production in η -quarkonium decays is carried out. The dependences of the studied characteristics on the mass of heavy quark and Higgs boson are investigated.

Yerevan Physics Institute

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Р.А.АЛАНАКЯН, С.Г.ГРИГОРЯН, С.Г.МАТИНЯН

РОЖДЕНИЕ ХИГГСОВСКОГО БОЗОНА В АДРОННЫХ РАСПАДАХ С-ЧЕТНЫХ КВАРКОНИЕВ

Изучены дифференциальные характеристики и полные ширины для рождения H^0- бозона в адронных распадах С-четных тяжелых кварконией (η , χ_0 , χ_1 , χ_2). Проведено сравнение дифференциальных распределений и полных ширин для рождения скалярного и псевдоскалярного бозонов в распадах η -кваркония. Исследованы зависимости изученных характеристик от массы тяжелого кварка и хиггсовского бозона.

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1. Introduction

$W^{\pm}$ and Z^0 bosons having been experimentally detected [1,2] , of peculiar interest became the search for the neutral scalar particle (H^0-boson) which is the last lacking element of the unified theory of Glashow-Salam-Weinberg (GSW) electroweak interaction [3-5]. Standard GSW model predicts the existence of single Higgs boson; however the search for the H^0-boson is complicated for its mass is not fixed in the model. The restrictions imposed on the mass are rather wide [6-8] :

$$7 \text{ GeV} \lesssim m_H \lesssim 1 \text{ TeV}$$

One should however mind that in the theories requiring the existence of two Higgs doublets*), the lower bound for the mass of one of them can be rather small.

Many recent works are devoted to searching for and studying the H^0-boson production mechanisms in various reactions [9-23] (pp, pp -collisions. e^+e^- -annihilation, vector quarkonium decays. etc.) with selection of the ones most advantageous as to their production cross-section magnitudes and

*) Note that supersymmetrization of standard model necessiates introduction of at least a second doublet of scalar bosons.

ways of their detection. Here the basic criterion of selecting certain mechanisms is that the interaction constant of Higgs particle is proportional to the mass of the particles it comes into interaction with. Besides, the improvement of the detection properties of a certain reaction (when, e.g. the H^0-boson is produced in association with the Z^0-boson, heavy quarks, hadronic jets, etc.) enables one to separate from the background phenomena which unfortunately take place and are large practically for any mechanisms of the Higgs particle production.

The decays of mesons containing heavy quarks are one of the best sources of information on production and properties of the Higgs bosons with mass $m_H < 2m$, where m is the mass of heavy quark. For vector quarkonia J/Ψ , Υ , T ($Q\bar{Q}$ states with $J^{PC} = 1^{--}$), their radiative decays with H^0-boson production were studied in the framework of nonrelativistic potential model [16] (Wilczek mechanism, see Fig.1). For these decays widths there was obtained (see also [14,15]):

$$\frac{\Gamma(V_{Q\bar{Q}} \to H^0 \gamma)}{\Gamma(V_{Q\bar{Q}} \to \mu^+\mu^-)} = \frac{G_F \, M_V^2}{4\sqrt{2}\,\pi\alpha} \left(1 - \frac{m_H^2}{M_V^2}\right) , \tag{1}$$

where $M_V \, (\simeq 2m)$ is the vector quarkonium mass. At $\frac{m_H^2}{M_V^2} \ll 1$ from (1) for the ratio of the quarkonium decay width by the Wilczek mechanism to the width of decay into muon pair we obtain (for M_V the following values are taken: $M_\Upsilon = 9.5$ GeV, $M_T = 80$ GeV):

$$\frac{\Gamma(\Upsilon \to H^0\gamma)}{\Gamma(\Upsilon \to \mu^+\mu^-)} \simeq 8 \;\, 10^{-3} , \qquad \frac{\Gamma(T \to H^0\gamma)}{\Gamma(T \to \mu^+\mu^-)} \simeq 0.56 ,$$

which already for bottonium will give a quantity quite accessible for experimental study. Note, however, that in Ref. 15 the ratio of widths for toponium was estimated without account of the Z^0-boson contribution to the width $\Upsilon \to \mu^+\mu^-$ which becomes noticeable for $M_\Upsilon \simeq 80$ GeV.

2. C-Even Quarkonia Decays with H^0- and P^0-Boson Production

In this paper we are proceeding farther with studying the mechanisms of H^0-boson production in the decays of C-even quarkonia such as $\eta(J^{PC}=0^{-+})$ [21], χ_J ($J^{PC} = 0^{++}, 1^{++}, 2^{++}$). Here the hadronic mechanism of Higgs boson production is the most advantageous one:

$$\eta(\chi_J) \to (H^0 + X), \tag{2}$$

where X stands for hadrons not containing heavy quarks.

At a quark-gluon level the width of decay (2) is determined by the decay (see Fig.2):

$$Q\bar{Q} \to H^0 + 2g, \tag{3}$$

where two gluons give finally the jets of ordinary hadrons, and Q and $\bar{Q}$ denote heavy quarks.

The amplitude of the process shown in Fig.2 has the following form:

$$M = 4\pi\alpha_s \cdot m\,(G_F\sqrt{2})^{1/2}\cdot\frac{1}{\sqrt{3}}\,Sp\left(\frac{\lambda^a}{2}\cdot\frac{\lambda^b}{2}\right)\cdot\varepsilon^a_\mu(K_1)\,\varepsilon^b_\nu(K_2)\,H(K_3)\times$$
$$\times\,\bar{v}(-K_5)\,\frac{\hat{K}_3-\hat{K}_5+m}{(K_3-K_5)^2-m^2}\,\gamma_\nu\,\frac{\hat{K}_4-\hat{K}_1+m}{(K_4-K_1)^2-m^2}\,\gamma_\mu\,u(K_4)+ \tag{4}$$

+ contributions of diagrams with permutations,

where $a, b = 1,\ldots, 8$, λ^a - are the Gell-Mann matrices of the SU(3) group; $\varepsilon^a_\mu(K_1)$, $\varepsilon^b_\nu(K_2)$ - are four-dimensional polarization vectors of gluons. $H(K_3)$ – is the H^0-boson wave function, α_s – is the strong interaction constant. The factor $\frac{1}{\sqrt{3}}$ is connected with color averaging. Notations for the particle momenta are given in Fig.2.

We are working within the nonrelativistic potential model [24], where the heavy quark momentum $\vec{P}$ in the c.m.s. ($\vec{K}_4 = -\vec{K}_5 \equiv \vec{P}$) is a natural parameter of expansion ($|\vec{P}| \ll m$), and since we are interested in decays of both S- and P-states, we have to extract from the amplitude expansion the zero- and first-order terms over momentum $\vec{P}$. A final expression for the considered amplitude in the Coulomb gauge is as follows:

$$
\begin{aligned}
M = & -4\sqrt{2}\,\pi\alpha_s \cdot m\left(G\sqrt{2}\right)^{1/2} \cdot \frac{\delta^{ab}}{2\sqrt{3}} H(K_3) \times \\
& \times \Big\{ i\,\frac{4m^2 - m_H^2}{2m^2(m_H^2 - 2m\varepsilon_3)} (n_2^i - n_1^i)\, e_1^j e_2^k \cdot \varepsilon_{ijk}\,\lambda_0 + \\
& + \frac{\varepsilon_3 - 2m}{m(m_H^2 - 2m\varepsilon_3)} \cdot \Big[\Big((\vec{n}_1\vec{e}_2)(\vec{\chi}\vec{e}_1) - (\vec{n}_1\vec{\chi})(\vec{e}_1\vec{e}_2) \Big) \times \\
& \times \Big((\vec{n}_1\vec{\xi}) + \frac{2m\varepsilon_3}{m_H^2 - 2m\varepsilon_3} (\vec{n}_3\vec{\xi}) \Big) - \frac{2m}{\varepsilon_1} (\vec{\chi}\vec{e}_2)(\vec{\xi}\vec{e}_1) + (1 \leftrightarrow 2) \Big] - \\
& - \frac{\varepsilon_3}{m(m_H^2 - 2m\varepsilon_3)} \Big[(\vec{\chi}\vec{\xi}) \cdot \Big((\vec{n}_1\vec{n}_3)(\vec{e}_1\vec{e}_2) - (\vec{n}_1\vec{e}_2)(\vec{n}_3\vec{e}_1) - (\vec{e}_1\vec{e}_2) \Big) + \\
& + \Big(-(\vec{n}_3\vec{\chi})(\vec{e}_1\vec{e}_2) + (\vec{n}_3\vec{e}_1)(\vec{\chi}\vec{e}_2) - (\vec{n}_3\vec{e}_2)(\vec{\chi}\vec{e}_1) \Big) \times \\
& \times \Big((\vec{n}_1\vec{\xi}) + \frac{2m\varepsilon_3}{m_H^2 - 2m\varepsilon_3} (\vec{n}_3\vec{\xi}) \Big) + (1 \leftrightarrow 2) \Big] +
\end{aligned}
\qquad (5)
$$

$$+\frac{1}{2m^2}(\vec{\chi}\vec{\xi})\cdot\left[(\vec{\ell}_1\vec{\ell}_2)-(\vec{n}_1\vec{n}_2)(\vec{\ell}_1\vec{\ell}_2)+(\vec{n}_1\vec{\ell}_2)(\vec{n}_2\vec{\ell}_1)\right]+$$

$$+\frac{1}{2m^2}\left[(\vec{n}_1\vec{\xi})-(\vec{n}_2\vec{\xi})\right]\cdot\left[(\vec{n}_1\vec{\ell}_2)(\vec{\chi}\vec{\ell}_1)-\right.$$

$$\left.-(\vec{n}_1\vec{\chi})(\vec{\ell}_1\vec{\ell}_2)-(\vec{n}_2\vec{\ell}_1)(\vec{\chi}\vec{\ell}_2)+(\vec{n}_2\vec{\chi})(\vec{\ell}_1\vec{\ell}_2)\right]- \tag{5}$$

$$\left.-\frac{1}{m\varepsilon_1\varepsilon_2}\cdot\left[\varepsilon_2(\vec{\xi}\vec{\ell}_1)(\vec{\chi}\vec{\ell}_2)+\varepsilon_1(\vec{\xi}\vec{\ell}_2)(\vec{\chi}\vec{\ell}_1)\right]\right\},$$

where $\vec{n}_\alpha = \frac{\vec{K}_\alpha}{\varepsilon_\alpha}$, α = 1,2,3 ; here ε_1 , ε_2-are gluon energies: ε_3 – is H^0-boson energy, ε_4 , ε_5 – are quark and antiquark energies in quarkonium; in the nonrelativistic limit in c.m.s. $\varepsilon_4 = \varepsilon_5 \simeq m$: $\vec{\ell}_1$ and $\vec{\ell}_2$ – are the gluon polarization vectors (in the Coulomb gauge $\varepsilon_0(K_1)=\varepsilon_0(K_2)=0$, $\varepsilon_i(K_1)\equiv \ell_{1i}$, $\varepsilon_i(K_2)\equiv \ell_{2i}$): $\vec{\xi} = \frac{\vec{P}}{m}$.

The Dirac spinors used are normalized by the condition $\bar{u}u$ = 1. Here $\sqrt{2}\chi_0 = Sp\,\chi$, $\sqrt{2}\vec{\chi} = Sp(\vec{\sigma}\chi)$, where $\chi^\alpha_\beta = \omega^\alpha \omega'_\beta$. ω^α and ω'^β are two-component spinors of quark and antiquark. The terms proportional to χ_0 correspond to spin 0 states (they actually correspond to contributions coming from η -quarkonium to the considered decay), while the ones proportional to $\vec{\chi}$ correspond to spin 1 states (to contributions coming from three χ_J -states to amplitude (4)).

Amplitude M is written so (see formula (5)). that the terms with χ_o will describe annihilation of η-quarkonium (^1S_o-state) to the system $2g\,H^0$ (two gluons and H^0-boson), whereas the terms with $\vec{\chi}$ - annihilation of χ_J-quarkonia (^3P_o, $^3P_1$, $^3P_2$ -states) - to the system $2g\,H^0$. To calculate the decay widths of χ_J-quarkonia, one should extract contributions from these three states to amplitude M. To do this, the most convenient way is as follows: we have at our disposal six independent vectors characterizing the decay: $\vec{\chi}$ (spin), $\vec{\xi}$, $\vec{\ell}_1$, $\vec{\ell}_2$, $\vec{n}_1$ and $\vec{n}_2$. From $\vec{\chi}$ and $\vec{\xi}$ we can construct a scalar ($\vec{\chi}\cdot\vec{\xi}$) corresponding to ^3P_o-state of quarkonium, a vector $[\vec{\chi}\,\vec{\xi}\,]$ corresponding to $^3P_1$-state, and a tensor $T_{\rho\sigma}=\chi_\rho\xi_\sigma+\chi_\sigma\xi_\rho - \frac{2}{3}(\vec{\chi}\,\vec{\xi})\,\delta_{\rho\sigma}$ corresponding to $^3P_2$-state. Together with the four vectors $\vec{\ell}_1$, $\vec{\ell}_2$, $\vec{n}_1$, $\vec{n}_2$ we must sort out all possible scalar combinations which will just determine the general structure of transition amplitude of the corresponding χ_J-quarkonia into the system of two gluons and H^0-boson. As a result we have:

For χ_o:

$$M_o = N(\vec{\chi}\,\vec{\xi})\cdot\left[R_1(\vec{\ell}_1\vec{\ell}_2) + R_2(\vec{n}_1\vec{\ell}_2)(\vec{n}_2\vec{\ell}_1)\right]; \qquad (6)$$

for χ_1:

$$\begin{aligned} M_1 = N[\vec{\chi}\,\vec{\xi}\,]\cdot\Big\{ & R_3[\vec{\ell}_1\vec{\ell}_2] + \\ & + R_4(\vec{\ell}_1\vec{\ell}_2)[\vec{n}_1\vec{n}_2] + R_5(\vec{n}_1\vec{\ell}_2)[\vec{n}_2\vec{\ell}_1] + \\ & + R_6(\vec{n}_2\vec{\ell}_1)[\vec{n}_1\vec{\ell}_2] + R_7(\vec{n}_1\vec{\ell}_2)[\vec{n}_1\vec{\ell}_1] + \\ & + R_8(\vec{n}_2\vec{\ell}_1)[\vec{n}_2\vec{\ell}_2]\Big\}, \end{aligned} \qquad (7)$$

for χ_2 :

$$M_2 = N T^{\rho\sigma} \cdot \left[R_9 \ell_1^\rho \ell_2^\sigma + R_{10}(\vec{n}_1 \vec{\ell}_2) \ell_1^\rho n_1^\sigma + \right.$$

$$+ R_{11}(\vec{n}_1 \vec{\ell}_2) \ell_1^\rho n_2^\sigma + R_{12}(\vec{n}_2 \vec{\ell}_1) \ell_2^\rho n_2^\sigma +$$

$$+ R_{13}(\vec{n}_2 \vec{\ell}_1) \ell_2^\rho n_1^\sigma + R_{14}(\vec{\ell}_1 \vec{\ell}_2) n_1^\rho n_1^\sigma + \tag{8}$$

$$\left. + R_{15}(\vec{\ell}_1 \vec{\ell}_2) n_2^\rho n_2^\sigma + R_{16}(\vec{\ell}_1 \vec{\ell}_2) n_1^\rho n_2^\sigma \right]$$

where the factor $N = -4\sqrt{2}\,\pi \alpha_s \cdot m (G\sqrt{2})^{1/2} \cdot \frac{\delta^{ab}}{2\sqrt{3}} \cdot H(K_3)$ is introduced into formulae (6), (7) and (8) for convenience. Everywhere the scalar product ($\vec{n}_1 \cdot \vec{n}_2$) is included to the coefficients R_i (i = 1, 2,..., 16). Comparing terms proportional to $\vec{x}$ in formula (5) with the sum $M_0 + M_1 + M_2$ we shall obtain for the coefficients R_i a system of equations which being solved will readily allow us to find

$$R_1 = \frac{2}{3}\left[R_9 + R_{14} + R_{15} + R_{16}(\vec{n}_1 \vec{n}_2) \right] + \frac{1-(\vec{n}_1 \vec{n}_2)}{2m^2} +$$

$$+ \frac{2\varepsilon_9}{m(m_H^2 - 2m\varepsilon_3)} \left[1 + \frac{2m-\varepsilon_3}{2\varepsilon_3} \left(1 + (\vec{n}_1 \vec{n}_2) \right) \right] , \tag{9}$$

$$R_2 = \frac{(4m^2 - m_H^2)^2}{6m^2(m_H^2 - 2m\varepsilon_3)^2} ,$$

$$R_3 = \frac{(\varepsilon_1 - \varepsilon_2)(4m^2 - m_H^2)}{2m\varepsilon_1\varepsilon_2(m_H^2 - 2m\varepsilon_3)} ,$$

$$R_4 = \frac{(\varepsilon_1 - \varepsilon_2)(4m^2 - m_H^2)}{2m(m_H^2 - 2m\varepsilon_3)^2} ,$$

$$R_5 = -R_{11} = \frac{(4m^2 - m_H^2)(2m(\varepsilon_3 - \varepsilon_2) - m_H^2)}{4m^2(m_H^2 - 2m\varepsilon_3)^2} , \tag{9}$$

$$R_6 = -R_{13} = R_5(\varepsilon_2 \longrightarrow \varepsilon_1) ,$$

$$R_7 = -R_{10} = \frac{(4m^2 - m_H^2)(m_H^2 - 2m(\varepsilon_1 + \varepsilon_3))}{4m^2(m_H^2 - 2m\varepsilon_3)^2} ,$$

$$R_8 = -R_{12} = R_7(\varepsilon_1 \longleftrightarrow \varepsilon_2) . ,$$

$$R_9 = \frac{(\varepsilon_1 + \varepsilon_2)(4m^2 - m_H^2)}{2m\varepsilon_1\varepsilon_2(m_H^2 - 2m\varepsilon_3)} ,$$

$$R_{14} = \frac{(m_H^2 - 2m(\varepsilon_1 + \varepsilon_3))(4m(m - \varepsilon_1) - m_H^2)}{4m^2(m_H^2 - 2m\varepsilon_3)^2} ,$$

$$R_{15} = R_{14}(\varepsilon_1 \longleftrightarrow \varepsilon_2) ,$$

$$R_{16} = \frac{8m^2\varepsilon_1\varepsilon_2 + (4m^2 - m_H^2)(m(2\varepsilon_3 - \varepsilon_1 - \varepsilon_2) - m_H^2)}{2m^2(m_H^2 - 2m\varepsilon_3)^2},$$

where $(\vec{n}_1 \vec{n}_2) = \frac{4m^2 - m_H^2 - 4m(\varepsilon_1 + \varepsilon_2) + 2\varepsilon_1\varepsilon_2}{2\varepsilon_1\varepsilon_2}$,

$$\varepsilon_3 = 2m - \varepsilon_1 - \varepsilon_2 \, .$$

Thus, formulae (5), (6), (7), (8) and (9) determine entirely the amplitudes of η and χ_J -quarkonia decay into $2g\ H^0$-system. To obtain from these the decay width Γ, one should take into account that (for details see Ref. [24])

$$\Gamma(i \to f) = \int |A(i \to f)|^2 d\Phi, \qquad (10)$$

where i is the quarkonium initial state, f - are the decay products, $d\Phi$ - is phase volume of final system, and the amplitude $A(i \to f)$ is related to our amplitudes M as follows:

$$A(i \to f) = \int M(\vec{p}) \cdot \varphi(\vec{p}) \frac{d^3p}{(2\pi)^3} \qquad (11)$$

here $\varphi(\vec{p})$ is quarkonium nonrelativistic wave function. For the considered C-even quarkonium states it has the following form:

$$\varphi(^1S_0) = \chi_0 \cdot \Psi_S(\vec{p}), \qquad (12)$$

$$\varphi({}^3P_0) = \frac{1}{\sqrt{3}}\left(\vec{\chi}\cdot\vec{\Psi}_p(\vec{p})\right),$$

$$\varphi_i({}^3P_1) = \frac{1}{\sqrt{6}}\,\varepsilon_{ijk}\cdot\chi_j\left(\Psi_p(\vec{p})\right)_k,$$

$$\varphi_{\rho\sigma}({}^3P_2) = \frac{1}{\sqrt{20}}\left(\chi_\rho(\Psi_p(\vec{p}))_\sigma + \chi_\sigma(\Psi_p(\vec{p}))_\rho - \frac{2}{3}(\vec{\chi}\vec{\Psi}_p(\vec{p}))\delta_{\rho\sigma}\right),$$

here $\Psi_s(\vec{p})$ and $\Psi_p(\vec{p})$ are Fourier images of $\Psi_s(\vec{r}) = \frac{R_s(r)}{\sqrt{4\pi}}$ and $\vec{\Psi}_p(\vec{r}) = \frac{\vec{r}}{r}\sqrt{\frac{3}{4\pi}}R_p(r)$ – wave functions of S and P -quarkonia, and $R_S(r)$ and $R_P(r)$ satisfy the normalization condition $\int|R(r)|^2 r^2 dr = 1$. The factors $\frac{1}{\sqrt{6}}$ and $\frac{1}{\sqrt{20}}$ take into account spin-averaging of ${}^3P_1$ - and ${}^3P_2$ -states of quarkonium.

Finally, for the widths of the studied decays we obtain:

$$d\Gamma(\eta\to 2gH^\circ) = \frac{2}{3\pi^2}\left(G_F\sqrt{2}\alpha_s^2\right)|R(0)|^2\times$$

$$\times\frac{(4m^2-m_H^2)^2\cdot(4m(\varepsilon_1+\varepsilon_2)+m_H^2-4m^2)^2}{2^6 m^2\varepsilon_1^2\varepsilon_2^2\,(m_H^2-2m\varepsilon_3)^2}\,d\varepsilon_1 d\varepsilon_2\,,$$

$$d\Gamma(\chi_0\to 2gH^\circ) = \frac{2}{3\pi^2}\left(G_F\sqrt{2}\alpha_s^2\right)|R_P'(0)|^2\times \tag{13}$$

$$\times\frac{9}{8}\Big[(1+(\vec{n}_1\vec{n}_2)^2)\cdot R_1^2 + (1-(\vec{n}_1\vec{n}_2)^2)^2\cdot R_2^2 -$$

$$-2(\vec{n}_1\vec{n}_2)(1-(\vec{n}_1\vec{n}_2)^2)R_1R_2\Big]d\varepsilon_1 d\varepsilon_2\,,$$

$$d\Gamma(\chi_1\to 2gH^\circ) = \frac{2}{3\pi^2}\left(G_F\sqrt{2}\alpha_s^2\right)|R_P'(0)|^2\times$$

$$\times \frac{1}{4}\Big\{2(\vec{n}_1\vec{n}_2)\big[(\vec{n}_1\vec{n}_2)^2-1\big]\cdot\big[R_4\cdot(-R_3+R_7-R_8+(\vec{n}_1\vec{n}_2)R_5-$$

$$-(\vec{n}_1\vec{n}_2)R_6)+R_5(R_8+R_3-2R_7+(\vec{n}_1\vec{n}_2)R_6)+$$

$$+R_6\cdot(-R_3+R_7-2R_8)\big]+$$

$$+2\cdot\big[(\vec{n}_1\vec{n}_2)^2-1\big]\cdot\big[2R_3R_7-2R_3R_8+R_7R_8-R_7^2-$$

$$-R_8^2\big]+\big[1-(\vec{n}_1\vec{n}_2)^4\big]\cdot\big[R_4^2+R_5^2+R_6^2\big]+$$

$$+\big[3-(\vec{n}_1\vec{n}_2)^2\big]\cdot R_3^2\Big\}d\varepsilon_1\, d\varepsilon_2 .$$

We do not present here the expression for the χ_2 quarkonium decay as it is rather cumbersome.

To obtain total widths and differential distributions (e.g. versus the energy of one of the gluons) one has to integrate the expressions, given in formulae (13), over ε_1 and ε_2 . The range of these variables is as follows:

$$\frac{4m^2-m_H^2}{4m}-\varepsilon_1 \leqslant \varepsilon_2 \leqslant \frac{m}{m-\varepsilon_1}\cdot\left(\frac{4m^2-m_H^2}{4m}-\varepsilon_1\right),$$

$$0 \leqslant \varepsilon_1 \leqslant \frac{4m^2-m_H^2}{4m} . \qquad (14)$$

Thus, integrating over ε_2 . we shall obtain differential distributions in ε_1 for the reactions under study (see Figs 3-9). In these figures we present the distribution function $F(X, m, m_H)$, i.e. the decay width divided by the factor $\frac{2}{3\pi^2}(G\sqrt{2}\alpha_s^2)\,|R_S(0)|^2\,\frac{4m^2-m_H^2}{4m}$ for η-quarkonium and by the factor $\frac{2}{3\pi^2}(G\sqrt{2}\alpha_s^2)\,|R_P'(0)|^2\cdot\frac{4m^2-m_H^2}{4m}$ for χ_J-quarkonia, respectively (the distributions are given as functions of quantity $X=\varepsilon_1/m\cdot(1-\frac{m_H^2}{4m^2})$, where X varies from 0 to 1, see formula (14)).

Before discussing the results obtained, let us take notice of the fact that in all the figures where we have plotted the distribution functions for the H^0-boson production from the η-quarkonium decay (Figs 3-6), we present also curves (dotted line) for the production of the neutral pseudoscalar particle P^0 in the η-quarkonium decays (it is known that the pseudoscalar Higgs boson can take place if two Higgs doublets instead of one are introduced into the usual scheme. Then one more, neutral scalar boson appears along with the usual scalar H^0-boson, and, generally speaking, two pseudoscalar bosons appear, one of which is neutral. Relatively heavy pseudoscalar mesons arise also in the technicolor models).

We have assumed that the neutral pseudoscalar boson interacts with the quarks with the same coupling constant as the scalar boson does [25,15]. For the P^0 production cross section in $\eta\to 2gP^0$ reaction we obtain

$$d\Gamma(\eta\to 2gP^0)=\frac{2}{3\pi^2}(G_F\sqrt{2}\alpha_s^2)|R(0)|^2\frac{1}{2^6m^2\varepsilon_1^2\varepsilon_2^2(m_H^2-2m\varepsilon_3)^2}\cdot$$

$$\cdot\left[m_H^4(4m(\varepsilon_1+\varepsilon_2)+m_H^2-4m^2)^2+\right. \qquad (15)$$

$$\left.+16m^4(4\varepsilon_1\varepsilon_2-4m(\varepsilon_1+\varepsilon_2)+4m^2-m_H^2)^2\right]d\varepsilon_1 d\varepsilon_2\,.$$

In the Wilczek mechanism (see Fig.1) the cross sections of H^0 and P^0 produced from the radiative decays of vector quarkonium coincide [25,15]. This means that to discriminate the emission of scalar particle from that of pseudoscalar one is impossible by studying in these decays only the total widths and angular distributions. Therefore the study of the P^0-production in C-even quarkonia decays seems attractive. It is evident already from formulae (13) and (15) that the differential widths for H^0 and P^0 produced from the η-quarkonium decays differ from each other (see also Appendix). As is seen from Figs 3-6, the gluon energy differential distribution is strongly different for H^0 and P^0 production and highly sensitive to their masses. Thus, at $m_H = 10$ GeV (see Fig.3) the production of P^0 pseudoscalar in the hadronic mechanism is suppressed more than by an order of magnitude as compared to the production of H^0 scalar, if its coupling with fermions is of the same order as that of H^0. With increasing m_H all the cross sections decrease; however the P^0 production cross section decreases slower than the H^0 one, therefore at $m_H \simeq 40$ GeV (see Fig.5) these cross sections become equal, while for large m_H the P^0 production cross section is dominant (see Fig.4: at $m_H = 75$ GeV the H^0-boson production cross section is by two orders suppressed). Fig.6 shows the differential distributions as functions of heavy quark mass. One can see that with increasing the quark mass (i.e. with increasing the quarkonium mass), these distributions grow up (this is actually due to the increase of the phase volume), but this dependence is weak (in our consideration everywhere under quarkonium system we shall imply toponium ($m_t = 40 \pm 10$ GeV [26]), which, as distinct from bottonium, will make it possible to discriminate the hadron jets coming from gluons).

In Figs 7-9 we present the distribution functions $F(x, m, m_H)$ of x for the decays from χ_J-quarkonia ($^3P_0$, $^3P_1$, $^3P_2$ – states).

For all m and m_H the values of these functions for χ_J -quarkonia satisfy $F_{\chi_0} > F_{\chi_2} > F_{\chi_1}$ condition. Fig.8 shows that with increasing m_H the distribution function for χ_0 decreases; a similar decrease is observed also for χ_1 and χ_2 quarkonia, however for the η -quarkonium decays this decrease is by two or three orders faster. These distributions are more sensitive to the quark mass. Thus, with increasing the quark mass from 30 to 50 GeV (see Fig.9) the distributions functions change by a factor of two or three. It is interesting that for χ_J -quarkonia, with increasing m (i.e. with increasing the phase volume) a decrease of differential cross sections takes place as distinct from what we had for η -quarkonium.

It is also interesting to investigate the differential distributions of the considered decays from invariant mass of produced gluon pair $\Delta^2 = (K_1 + K_2)^2$ (more exact, from $z = \Delta^2/(4m^2 - m_H^2)$). To obtain differential distributions in z, it is sufficient to make replacements in formulae (13) and (15):

$$\varepsilon_2 = -\varepsilon_1 + \frac{4m^2 - m_H^2}{4m}(1+z),$$

$$\varepsilon_3 = \frac{4m^2 + m_H^2 - (4m^2 - m_H^2)\cdot z}{4m}, \qquad (16)$$

$$d\varepsilon_1 d\varepsilon_2 = \frac{4m^2 - m_H^2}{4m}\, dz\, d\varepsilon_1,$$

where the variation ranges over z and ε_1 are as follows:

$$\varepsilon_1^{min} \leqslant \varepsilon_1 \leqslant \varepsilon_1^{max} \; ,$$

$$0 \leqslant z \leqslant \frac{2m - m_H}{2m + m_H} \; , \qquad (17)$$

here

$$\varepsilon_1{}^{max}_{min} = \frac{1}{8m} \Big\{ (4m^2 - m_H^2)(1+z) \pm \big[(4m^2 - m_H^2)^2 \cdot (1+z)^2 -$$

$$- 16 m^2 (4m^2 - m_H^2) \cdot z \big]^{1/2} \Big\} \; .$$

Integrating further over ε_1 , we shall obtain differential distributions $R(z, m, m_H)$ as functions of z , presented in Figs 10-13. The functions $R(z, m, m_H)$ are again divided by factors $\frac{2}{3\pi^2}(G\sqrt{2}\alpha_S^2)\,|R_S(0)|^2$ for η -quarkonium and $\frac{2}{3\pi^2}(G\sqrt{2}\alpha_S^2)\,|R_P'(0)|^2$ for χ_J -quarkonia, respectively. One can see from the figures that the distributions in z are characterized by a peculiar behaviour. For example, whereas for the decay $\eta \to 2gH^\circ$ the distribution at $z \to 0$ tends to zero, for the decay $\eta \to 2gP^\circ$, irrespective of m and m_H , the distribution function at $z \to 0$ tends to 1. Besides, the distribution function $R(z, m, m_H)$ for the decays of both ${}^1S_0$ -state (η -quarkonium) and ${}^3P_J$ -states (χ_J -quarkonium) at $z \to z^{max}$ tends to zero (see Figs 10-13). Note that variation of these distributions with m and m_H is similar to the one we had for differential characteristics over X .

3. Background Processes.

It is thus obvious that the considered differential distributions $\frac{d\Gamma}{dx}$ and $\frac{d\Gamma}{dz}$ (for the Higgs boson production in hadronic decays of C-even quarkonia) have a characteristic behaviour and they are sensitive to mass parameters variation. These very properties may play an important role in searching for H^0-bosons on those decays and in their separation from the background processes. Thus, for the considered masses of H^0-boson ($m_H < 2m_t = 80$ GeV), its main decay channels are $H^0 \to \tau\bar{\tau}$, $H^0 \to c\bar{c}$, $H^0 \to b\bar{b}$. For the decay process $\eta(\chi_J) \to H^0 +$ two jets $\to \tau\bar{\tau}$ + two jets, i.e. when at the end we follow a pair of heavy leptons and two jets of light hadrons, the main background comes from vector quarkonia decays with lepton pair production via a virtual photon (see Fig.14). However, the differential distribution over x for this decay (Ore-Powell formula [27]) grows monotonously from 0 at $x = 0$ to 1 at $x = 1$, as distinct from the distribution, say, for η-quarkonium decay (see Fig.3), where in the region of large x we have a clearly pronounced maximum, while at $x = 1$ the distribution function turns to zero. Note also that this background can be easily separated since the quarkonium initial state is C-even (vector quarkonium). The background situation is somewhat worse for the decays of the type of $\eta(\chi_J) \to H^0$ + two jets $\to c\bar{c}$ or $b\bar{b}$ + two jets, where the main backgrounds come from decays of quarkonia (both C-odd and C-even) into three gluons with subsequent transition of one gluon to a quark pair $c\bar{c}$ or $b\bar{b}$ (the other two gluons produce two jets). Separation from these backgrounds is already a more complicated problem.
It should be noted that the presence of three-particle states in the final state of investigated decays (H^0-boson and two jets emerging only from decays of heavy quarkonium C-even states) gives rise to utilization of the new

characteristics of the differential distribution type, very sensitive, as we have seen (see also [21]), to the coupling of Higgs boson to fermions.

The considered differential characteristics of quarkonia hadronic decays with H^0-boson production contain infrared divergences. It is of interest that we have infrared divergences in distribution functions already for decays from S-states (see formula (A.1) of Appendix: when $m_H = 0$, i.e. $\lambda = 1$, at $x \to 1$ in the function $f(x, \lambda)$ we have a divergence due to the "softness" of H^0-boson).

In electrodynamics, the infrared photon production amplitude is proportional to the expression [28-30]:

$$\frac{(K_4 \cdot \varepsilon^*)}{(K_4 \cdot K)} - \frac{(K_5 \cdot \varepsilon^*)}{(K_5 \cdot K)} ,$$

where ε_μ is the infrared photon polarization vector, and K – is its 4-momentum. In quarkonium, this factor is replaced by $\frac{2(\vec{P} \cdot \vec{\varepsilon}^*)}{m\omega}$ [30], where ω is the soft photon energy. From here it is obvious that the soft photon always takes the orbital momentum away (since $\vec{P}$ – is the orbital motion momentum). That is infrared divergence emerges in P-wave only. We had this divergence already in S-wave (i.e. already in decays of η -quarkonium), this being due to a scalar nature of the H^0-boson-quark coupling. However due to $m_H \neq 0$ these divergences remove in a natural way. As to the P-state decays (χ_J -quarkonium), we, as well as in electrodynamics, have infrared divergences due to the soft gluons (see Figs 7-9, where these divergences manifest themselves at $\varepsilon_1 \to 0$ and $\varepsilon_1 \to 1$). Energy cut-off for these gluons is performed in a natural way, since in computer calculations these energies do not reach a zero value. Besides, in computing, the energy progressing from $\varepsilon_1 \simeq 100$ MeV even by two orders towards zero value caused variation in magnitudes of distribution functions and total widths no more than 10 ÷ 20 %, the latter being of the same order as uncertainties introduced by inaccuracies of QCD calculations and computer errors.

4. Total Widths

To obtain total widths of the studied decays, one should integrate the expressions in formulae (13) and (15) over ε_1 and ε_2. It is interesting to compare the obtained widths to the width of the vector quarkonium decay into $H^0\gamma$. Thus, for the η -quarkonium decay we obtain [21]

$$\frac{\Gamma(Q\bar{Q}(^1S_0) \to 2gH^0)}{\Gamma(V_{Q\bar{Q}} \to H^0\gamma)} = \frac{4\alpha_s^2}{3\pi\alpha Q_q^2} \cdot \int \left(mF(x,m,m_H)\right)dx \,, \tag{18}$$

where Q_q is heavy quark charge.

The values for this ratio are listed in Table 1.

Table 1

m_H , GeV	1	10	20	30	40	75
$\frac{\Gamma(\eta \to 2gH^0)}{\Gamma(V_{Q\bar{Q}} \to H^0\gamma)}$	4.0	1.5	0.6	0.3	0.15	1.2 10^{-3}

One can see that for not large masses of Higgs boson ($m_H \lesssim$ 10 GeV) the H^0-boson is produced in hadronic decays more frequently than in quarkonium radiative decays. These widths are comparable for m_H in the range 10 ÷ 30 GeV. With further increase of H^0-boson mass, this ratio sharply falls off reaching some fractions of a percent of H^0-boson production widths in radiative decays (for estimations, $\alpha_s(2m_t)$ is taken $\simeq$ 0.1).

It is of interest to investigate a ratio of widths of decays into the system $2g\,H^0$ from S- and P-states. Thus, in case of S- and P-state decays into two gluons one can obtain [31] :

$$\frac{\Gamma(^3P_0 \to 2g)}{\Gamma(^1S_0 \to 2g)} = 9 \cdot \left| \frac{R'_P(0)}{m R_S(0)} \right|^2 , \qquad (19)$$

this giving $\sim 10^{-2}$ for the Coulomb potential (a number close to that turns out to be for the Richardson potential as well [32]). In our case one can obtain

$$\frac{\Gamma(^3P_0 \to 2gH^0)}{\Gamma(^1S_0 \to 2gH^0)} = æ \left| \frac{R'_P(0)}{m R_S(0)} \right|^2 , \qquad (20)$$

where æ is strongly dependent on the Higgs boson mass. The values of æ for different m_H are listed in Table 2 (the quark mass is m = 40 GeV).

Table 2

m_H , GeV	1	10	20	40	75
æ	16.0	31.6	54.6	$1.7\ 10^2$	$1.7\ 10^4$

It is seen that the P-wave relative contribution increases in the presence of H^0-boson.

Thus, at m_H= 40 GeV the ratio of contributions of P- and S-state decay widths amounts to $\simeq$ 20%.

Hence one can see that the considered mechanism of Higgs boson production is interesting as to behaviour of both the differential characteristics (see Figs 3-13) and the total decay width (see formulae (19) and (20) and

Tables 1 and 2). Besides, the presence of two hadron jets together with H^0-boson in the final state also can help with identification of the Higgs particles. Note that C-even quarkonium states, that are necessary for the Higgs scalar production in the investigated decay, one can obtain either in hadron collisions (via the gluon-gluon fusion mechanism) or in e^+e^- - annihilation, where one can get C-even quarkonium states from radiative decays of vector quarkonia (in this case decays in χ_J-state proceed with large branching).

We present also the ratio of total widths of η -quarkonia decays into H^0-boson and P^0 pseudoscalar (Table 3) for different values of masses m_H.

Table 3

m_H , GeV	1	10	20	30	40	56.5	75
$\frac{\Gamma(\eta \to 2gH^0)}{\Gamma(\eta \to 2gP^0)}$	30.2	9.4	4.6	2.2	1.2	0.3	$8.6 \cdot 10^{-3}$

5. About One Mechanism of H^0-Boson Associated Production.

Let us finally discuss one more mechanism of H^0-boson associated production in quarkonium decays. In Fig.15 we give a diagram of H^0-boson production by Z^0-boson due to three-boson couplings (see, e.g. [11.15]). The production of Higgs particle in association with gauge Z^0-bosons creates conditions suitable enough for their identification. Note that this mechanism can proceed both from the vector quarkonium decay (due to the vector part of interaction of Z^0-boson with heavy quarks) and from the C-even η -quarkonium decay (due to the axial part of Z^0-boson interaction with heavy quarks).

Standard calculations for the widths of these decays lead to:

$$\frac{\Gamma(V_{Q\bar{Q}} \to H^o f\bar{f})}{\Gamma(V_{Q\bar{Q}} \to H^o \gamma)} = \frac{G_F^2}{4\pi^3 \alpha Q_q^2} \cdot \frac{V_t^2(V_f^2+a_f^2)\, m^2 m_Z^8}{(4m^2-m_H^2)^2 \cdot (4m^2-m_Z^2)^2} \cdot J_1(m_H), \tag{21}$$

$$\frac{\Gamma(\eta \to H^o f\bar{f})}{\Gamma(V_{Q\bar{Q}} \to H^o \gamma)} = \frac{3G_F^2}{4\pi^3 \alpha Q_q^2} \cdot \frac{a_t^2(V_f^2+a_f^2)\, m^2 m_Z^4}{4m^2-m_H^2} \cdot J_2(m_H),$$

here m_Z is the Z^0-boson mass; $f, \bar{f}$ - are quarks and leptons in final state (their masses are neglected as compared to m_t and m_Z); V_t , a_t , V_f and a_f are vector and axial constants of interaction of Z^0-boson with quarks and leptons; $V_t = 1 - \frac{8}{3}\sin^2\theta_W$, $a_t = 1$, $V_f = V_t$, $a_f = a_t$ for u and c quarks and $V_f = -1 + \frac{4}{3}\sin^2\theta_W$, $a_f = -1$ for d , s and b quarks; $V_f = -1 + 4\sin^2\theta_W$, $a_f = -1$ for electron and muon. From (21) and Table 4 one can see that the width of the considered decay is suppressed as compared to the Wilczek mechanism.

Table 4

m_H , GeV	1	10	40	m_H , GeV	1	10	40
$J_1(m_H)$	0.21	0.1	$4.6 \cdot 10^{-3}$	$J_2(m_H)$	0.01	$8.5 \cdot 10^{-3}$	$9.2 \cdot 10^{-4}$

Thus, for the vector quarkonium decay, this ratio for m_H= 10 GeV is $\sim 10^{-3}$. This ratio cannot be improved significantly if with respect to the quarkonium mass we shall be in the resonance region ($2m \simeq m_Z$). In case of η-quarkonium, this ratio is even lesser. Thus it is obvious that the considered mechanism of associated H^0-boson production via three-boson couplings HZZ in quarkonium decays is small.

Note also that we have examined also the heavy quarkonium decay into two standard H^0-bosons (see Fig.16 and also Ref. 23). Since the system $2H^0$ is C-even, then, with account of P-parity conservation, this system is produced only through χ_0- and χ_2-quarkonia decays.

Investigation of the mentioned decay is useful for extraction and study of a contribution of three-boson couplings that play such an important role in theory (HHH couplings in the given case).

Note that the HHH coupling contributes only to the decay proceeding from the χ_0-quarkonium. Calculations for the decay $\chi_0 \to 2H^0$ have shown that some suppression of this decay as compared to the Wilczek mechanism is due to the fact that the considered decay proceeds through the P-wave. So, for the widths ratio of these decays we obtain:

$$\frac{\Gamma(^3P_0 \to 2H^0)}{\Gamma(V_{Q\bar{Q}} \to H\gamma)} \simeq (10 \div 30) \cdot \left| \frac{R'_P(0)}{mR(0)} \right|^2 ,$$

i.e., e.g. for the Coulomb potential this ratio is $\sim (1 \div 3) \cdot 10^{-2}$ (for $m_H = 10 \div 30$ GeV and $m = 40$ GeV). Despite the suppression of $\chi_0 \to 2H^0$ decay, i.e. difficulty of its experimental detection, the study of this reaction undoubtedly provides an interesting and important information on the Higgs boson self-interaction.

Note that the $\chi_2 \to 2H^0$ decay turned out by two orders of magnitude more suppressed as against the $\chi_0 \to 2H^0$ one.

6. Conclusion

Using the results obtained, let us estimate the number of events for the H^0-boson production. Thus, in $p\bar{p}$ (pp) - collisions, for the production cross section of the Higgs boson decaying into the channel H^0 + two jets: $p\bar{p}$ (pp) $\to \eta_t + X \to H^0$ + two jets $+ X$, we have $\sigma(H^0) = \sigma(2g \to \eta_t)$

$Br\ (\eta_t \to 2gH^\circ)$, this giving for $Sp\bar{p}S$ -collider $\sigma(H^\circ) \sim$ $\sim (3 \div 0.2) \cdot 10^{-37}$ cm^2 at $m_H = (10 \div 40)$ GeV. For accessible luminosities of a collider ($L \simeq 2 \cdot 10^{30}$ cm^{-2}sec^{-1}), this provides production of 25 events per year for the H^0-boson production with $m_H = 10$ GeV. For Tevatron energies, the η_t -quarkonium production cross section is by one order higher, so we shall obtain several H^0-bosons per week. The possibility of the investigation of the large-mass Higgs boson will necessiate a substantial increase of luminocity on $Sp\bar{p}S$ -collider.

In e^+e^- -annihilation for the H^0-boson production cross section by the studied mechanism ($e^+e^- \to \Upsilon \to \chi_J \cdot \gamma \to H^\circ + \gamma$ + two jets), taking into account that $Br\ (\chi_0 \to 2gH^\circ) \sim (0.8 \div 0.4) \cdot 10^{-2}$, for $m_H = (10 \div 40)$ GeV, we have $\sigma(H^\circ) \sim (0.8 \div 0.4) \cdot 10^{-37}$ cm^2. This corresponds to one event per week for luminocities being planned at LEP.

In conclusion, the authors express their thanks to I.G. Aznauryan for the useful and interesting discussions, to A.Yu. Khodjamiryan for a series of valuable remarks, and to L.S. Dulyan for assistance in computing.

APPENDIX

Integrating expressions (13) and (15) over ε_2 , for η -quarkonium decays we obtain

$$d\Gamma(\eta \to 2gH^{\circ}) = \frac{2}{3\pi^2}(G_F\sqrt{2}\alpha_s^2)|R(0)|^2 \cdot f(x,\lambda)\,dx\,,$$

$$d\Gamma(\eta \to 2gP^{\circ}) = \frac{2}{3\pi^2}(G_F\sqrt{2}\alpha_s^2)|R(0)|^2 \times \tag{A.1}$$

$$\times\left\{(1-\lambda)^2+\lambda^2-2\lambda x+\lambda x^2\right\}\cdot f(x,\lambda)\,dx\,,$$

where

$$f(x,\lambda) = \frac{1-x}{\lambda x(2-x)^2}\cdot\left[\frac{2\lambda-2x+x^2}{\lambda-2x+x^2} + \right.$$

$$\left. + \frac{2\lambda}{x(2-x)}\cdot \ln\left(\frac{\lambda-2x+x^2}{\lambda}\right)\right],$$

$$x = \frac{\varepsilon_1}{m\left(1-\frac{m_H^2}{4m^2}\right)}\,, \quad \lambda = \frac{4m^2}{4m^2-m_H^2}\,.$$

From (A.1) it is seen that the distribution functions for H^0 and P^0 differ by factor $\left\{(1-\lambda)^2+\lambda^2-2\lambda x+\lambda x^2\right\}$ which is just responsible for the essential differences in their values and behaviours mentioned in the text (see Figs 3-6).

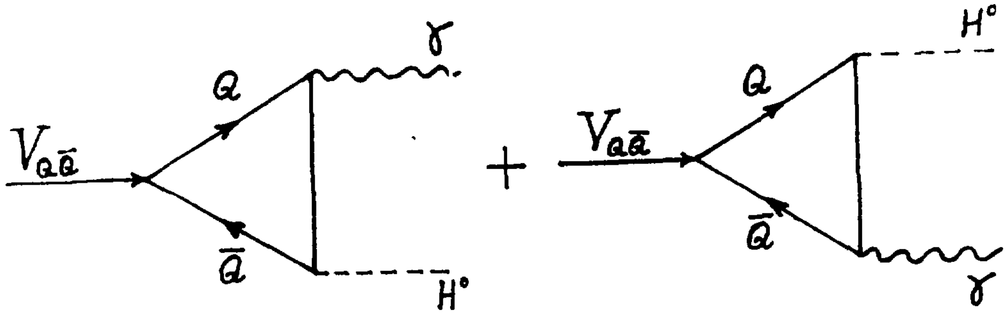


Fig.1. Diagrams describing the vector quarkonium radiative decay with H^0-boson production.

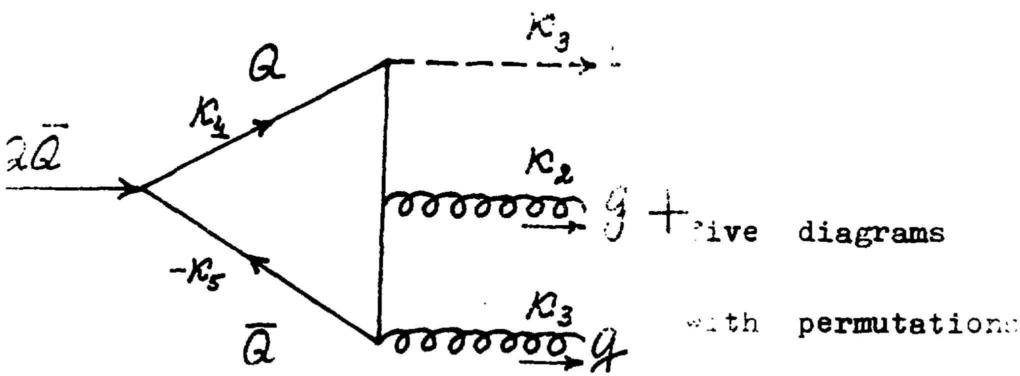
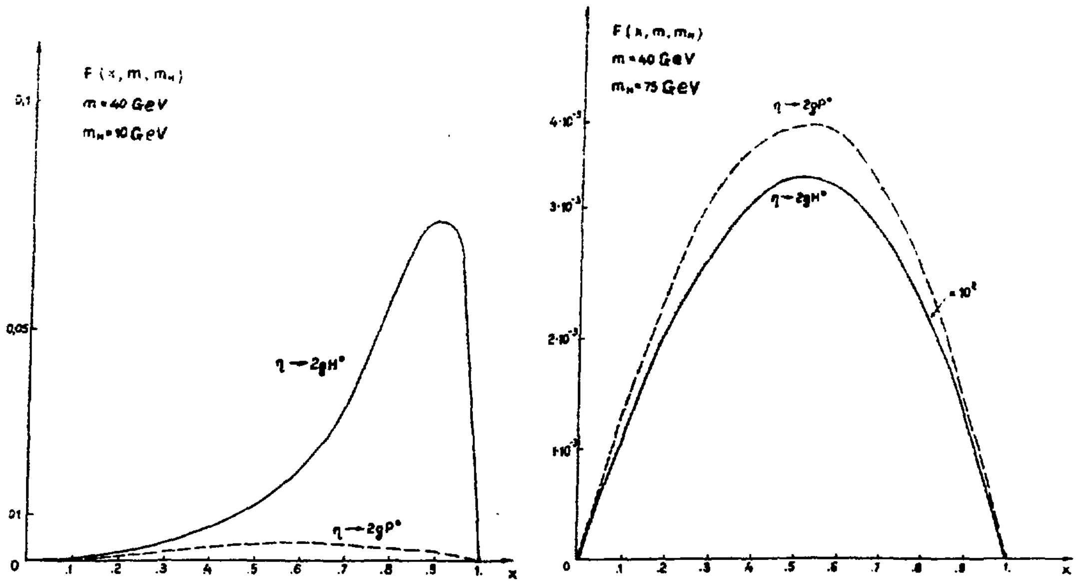


Fig.2. Diagrams corresponding to hadronic mechanism of H^0-boson production in heavy quarkonium decays.



Figs 3 and 4. Differential distribution $F(x, m, m_H)$ as a functio of $x = \varepsilon_1/m\left(1 - \frac{m_H^2}{4m^2}\right)$ for H^0- and P^0-boson production in η-quarkonium decays. Curves in Fig.3 are plotted for m = 40 GeV and m_H = 10 GeV; in Fig.4 for m = 40 GeV and m_H = 75 GeV.

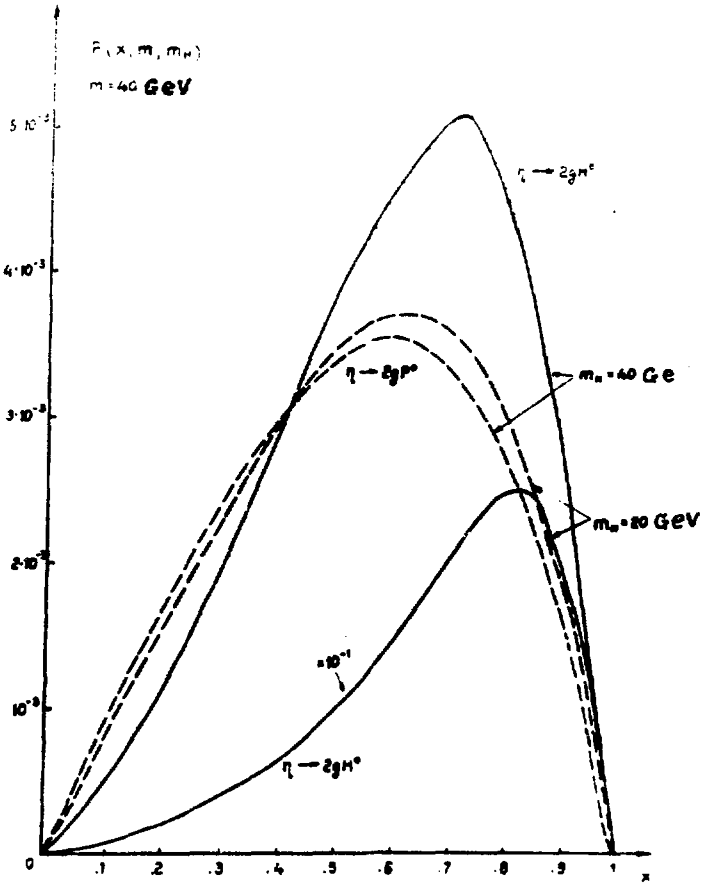
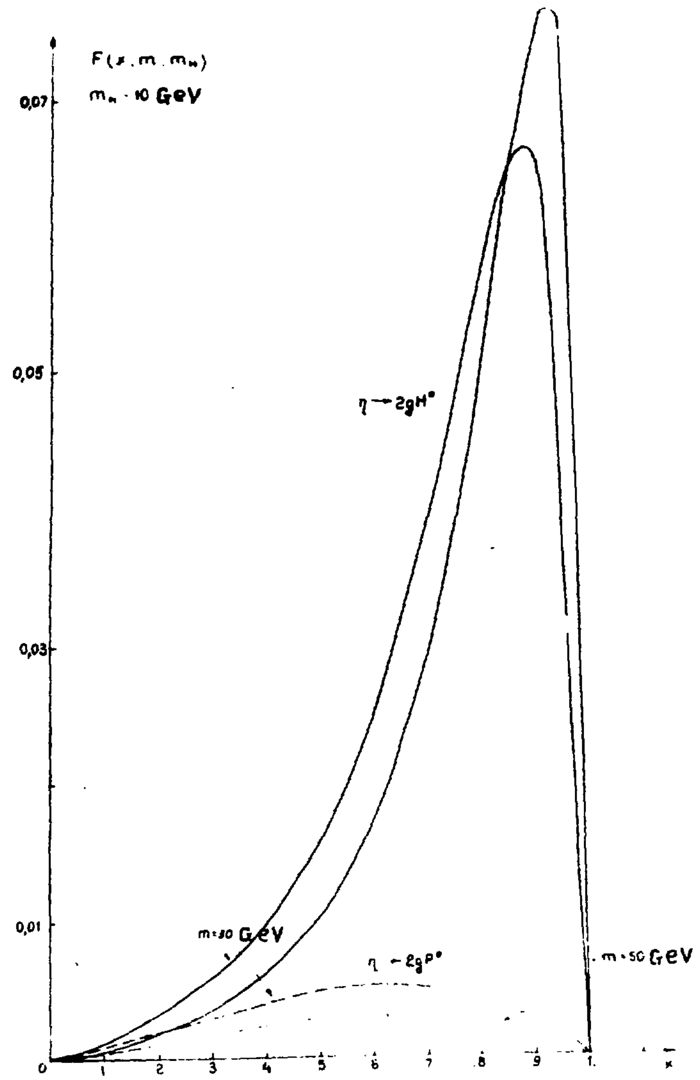


Fig.5. Dependence of $F(x, m, m_H)$ for H^0- and P^0-boson productic in η-quarkonium decays on the Higgs particle mass. Curves ar plotted for two values: m_H = 20 and 40 GeV (everywhere m = 40 GeV).



ig.6. Variation of differential distribution $F(x, m, m_H)$ for H^0- and P^0-boson production in η-quarkonia decays versus the heavy quark mass. Curves are plotted for heavy quark masses m = 30 and 50 GeV (H^0-boson mass is m_H = 10 GeV everywhere).

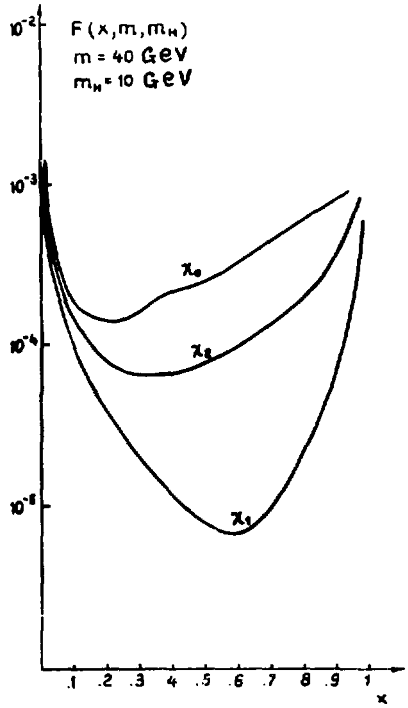


Fig.7. Differential distribution $F(x, m, m_H)$ as a function of x for H^0-boson production in χ_J -quarkonia decays. Curves are plotted for m = 40 GeV and m_H = 10 GeV.

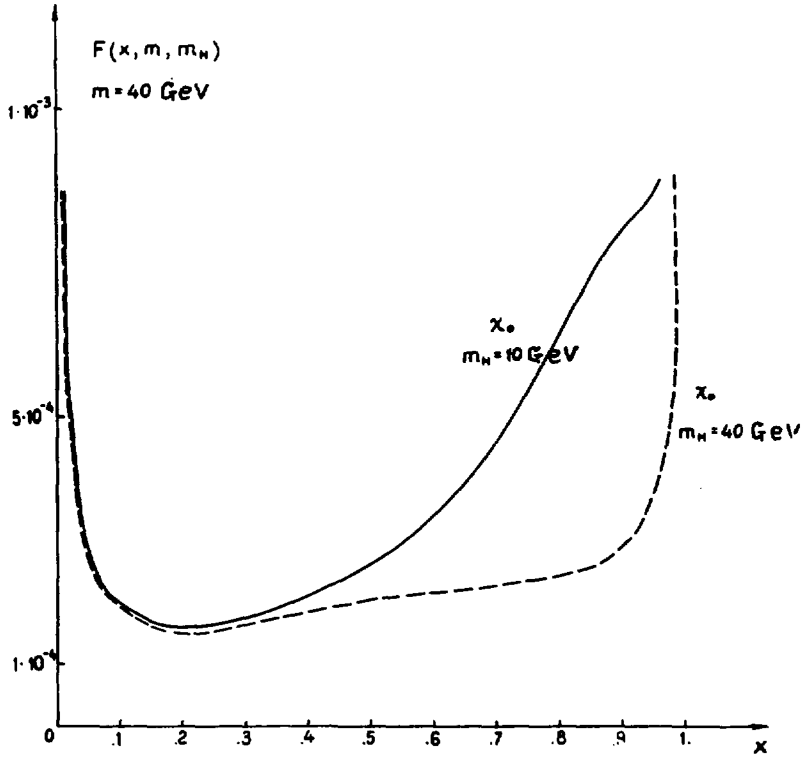


Fig.8. Dependence of differential distribution $F(x, m, m_H)$ for H^0-boson production in χ_0-quarkonium decay on the Higgs particle mass. Curves are plotted for m_H = 10 and 40 GeV and m = 40 GeV.

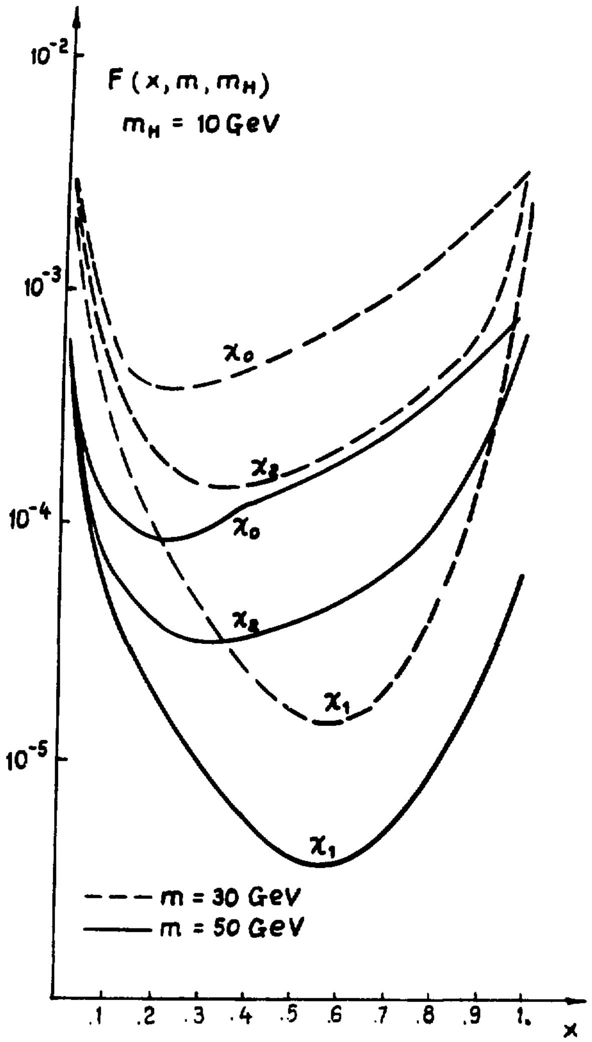
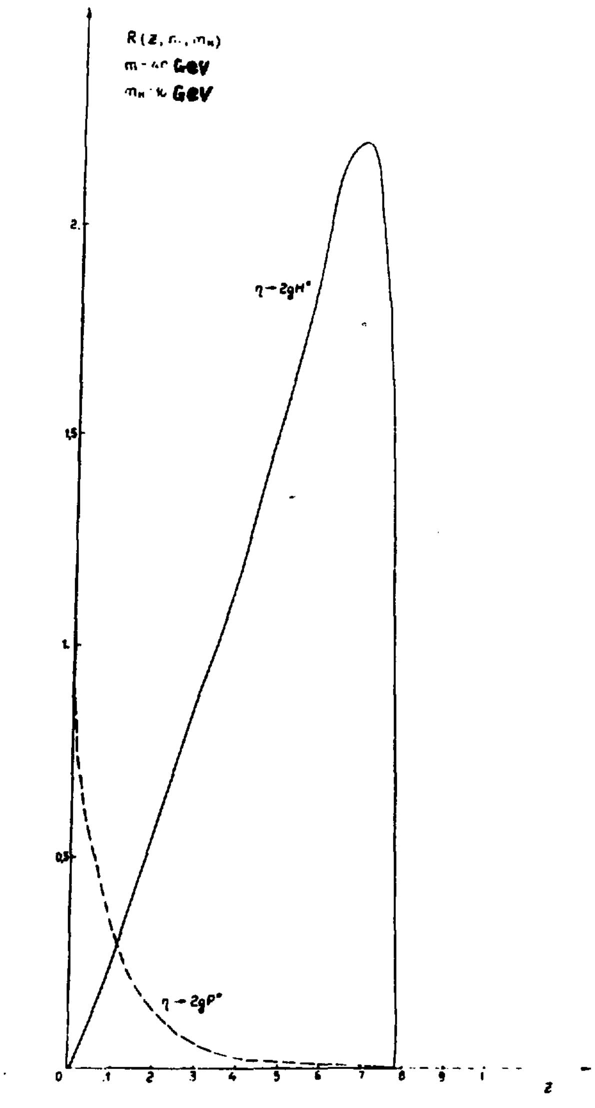
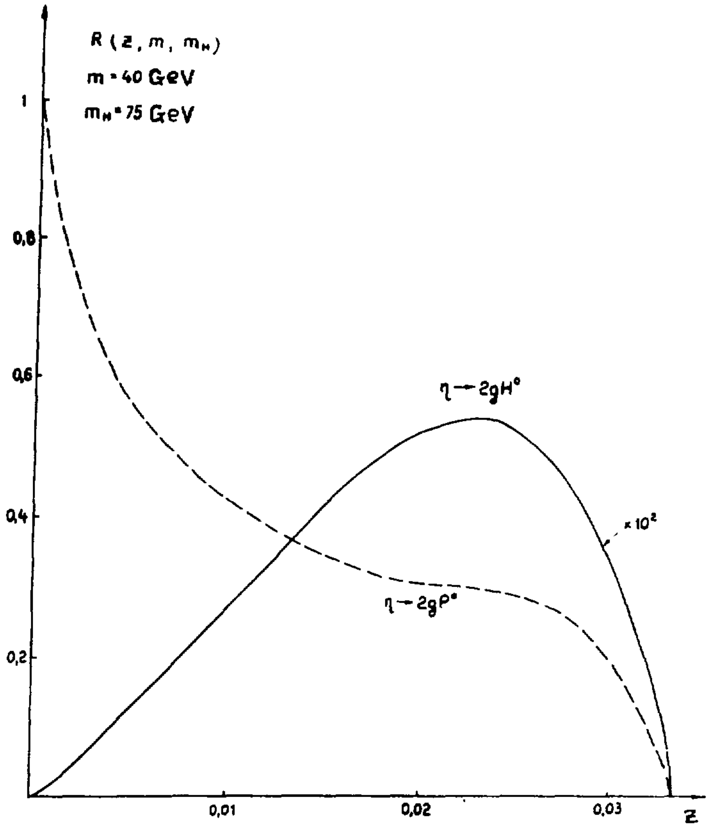


Fig.9. Differential distribution $F(x, m, m_H)$ for H^0-boson production in χ_J -quarkonia decays as a function of heavy quark mass. Curves are plotted for m = 30 and 50 GeV and m_H = 10 GeV.


GeV
GeV
η → 2gH°
η → 2gP°
2.
1.5
1.
0.5
0
z



Figs 10 and 11. Differential distribution $R(z, m, m_H)$ as a function of $z = \Delta^2/(4m^2 - m_H^2)$ for H^0 and P^0-boson production in η -quarkonium decays. Curves in Fig.10 are plotted for m = 40 GeV, m_H = 10 GeV; in Fig.11 for m = 40 GeV and m_H = 75 GeV.

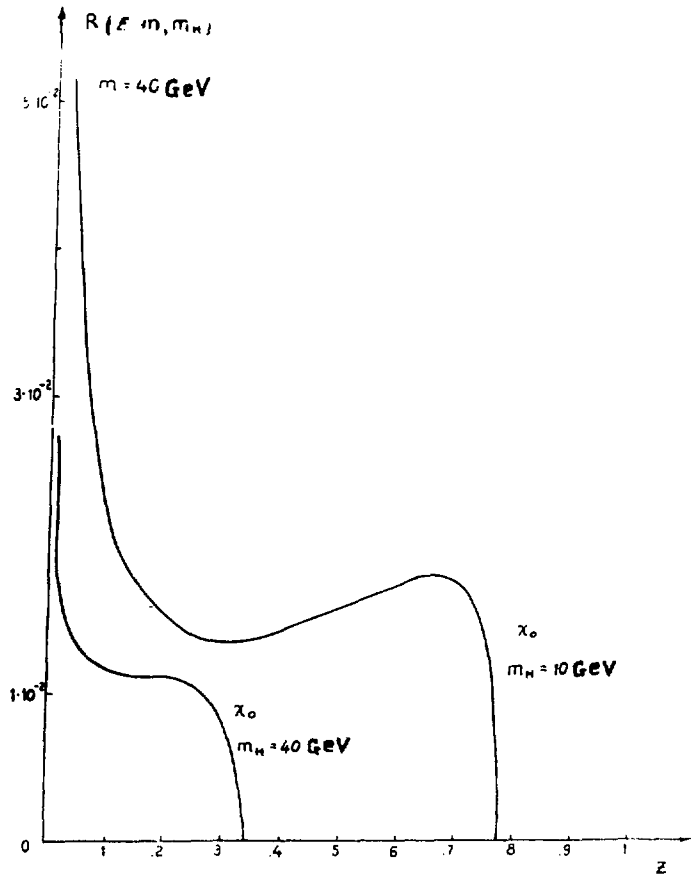


Fig.12. Differential distribution $R(z, m, m_H)$ for H^0-boson production in χ_o-quarkonium decay as a function of Higgs boson mass. Curves are plotted for m_H = 10 and 40 GeV and m = 40 GeV.

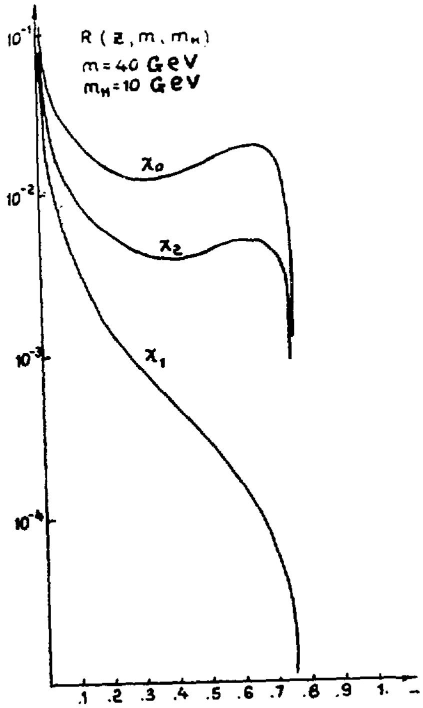
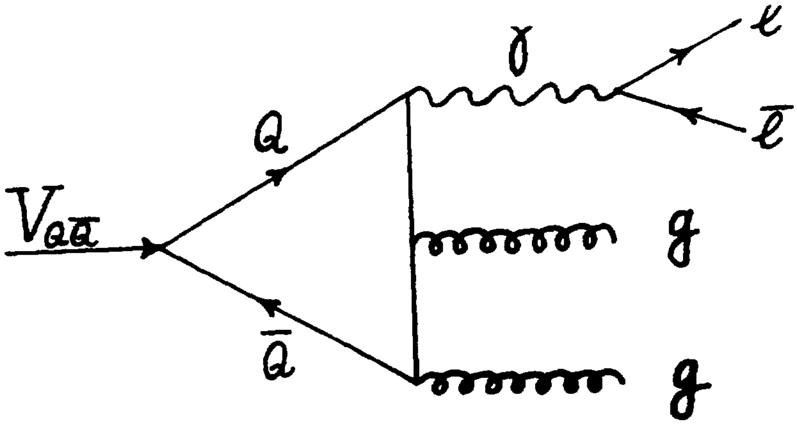


Fig.13. Differential distribution $R(z, m, m_H)$ as a function of z for H^0-boson production in χ_J -quarkonia decays. Curves are plotted for m = 40 GeV and m_H = 10 GeV.



ig.14. Diagram describing the vector quarkonium radiative decay with two-gluon production.

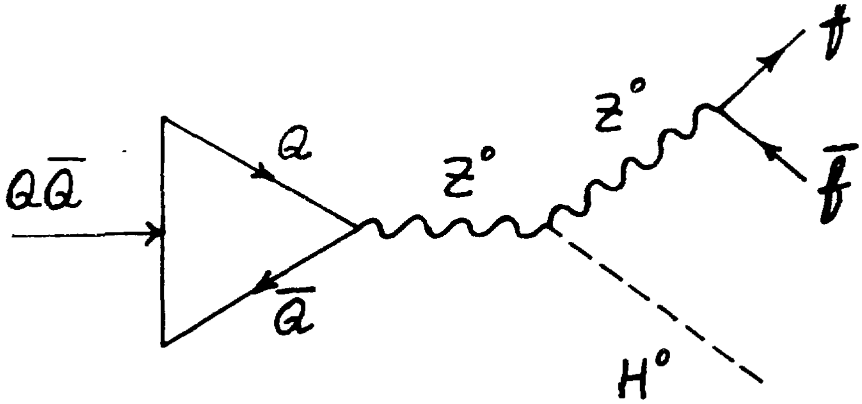


Fig.15. Diagram describing three-boson production of H^0-boson accompanied with Z^0-boson in heavy quarkonium decays.

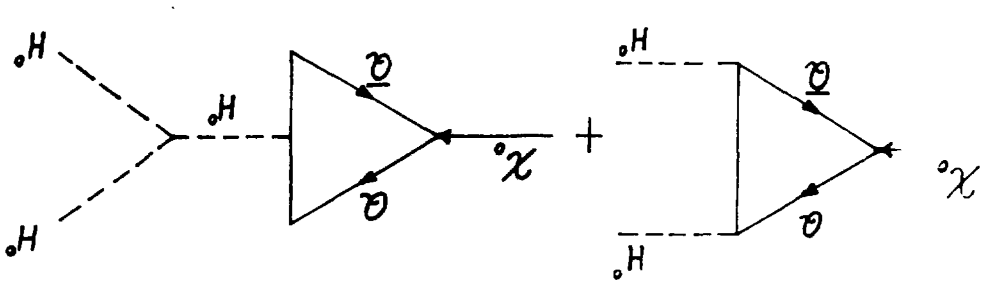


Fig.16. Diagrams describing the production mechanism of two standard H^0-bosons in C-even quarkonia decays.

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РОЖДЕНИЕ ХИГГСОВСКОГО БОЗОНА В АДРОННЫХ РАСПАДАХ С-ЧЕТНЫХ КВАРКОНИЕВ

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