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HEAVY QUARKONIA PRODUCTION IN HADRON-HADRON COLLISIONS

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ԾԱՆՐ ՔՎԱՐԿՈՆԻՈՒՄՆԵՐԻ ԾՆՈՒՄԸ ՀԱԴՐՈՆ - ՀԱԴՐՈՆԱՅԻՆ $^1P_1$ ՀԱՐՎԱԾՆԵՐՈՒՄ

Հետազոտված է մեծ լայնական իմպուլսներով $^1P_1$ քվարկոնիումների /չարմոնիումի, բոտոնիումի և տոպոնիումի/ ծնումը գերբարձր էներգիաներում pp և p̄p փնջերի բախումներում։ Քննարկումը կատարվում է քվարկ-գլյուոնային միավորման և ոչ ռելյատիվիստական քվարկոնիումների մոդելի շրջանակներում։

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РОЖДЕНИЕ $^1P_1$ ТЯЖЕЛЫХ КВАРКОНИЕВ В АДРОН-АДРОННЫХ СТОЛКНОВЕНИЯХ

Исследовано рождение $^1P_1$ - кваркониев (чармония, боттомония и топония) с большими поперечными импульсами в столкновениях pp и $p\bar{p}$ пучков сверхвысоких энергий. Рассмотрение проводилось в рамках механизма глюон-глюонного слияния и нерелятивистской модели кваркониев.

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$^1P_1$ HEAVY QUARKONIA PRODUCTION IN HADRON-HADRON COLLISIONS

The production of $^1P_1$ -quarkonia (charmonium, bottomonium and toponium) with large transverse momenta is studied in collisions of ultrahigh-energy pp and $p\bar{p}$ beams. The analysis was carried out in the framework of gluon-gluon fusion mechanism and nonrelativistic quarkonium model.

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The heavy quarkonium physics occupies a particular place in quantum chromodynamics (QCD). From the QCD viewpoint, the heavy quark $Q\bar{Q}$ pairproduction in hadronic collisions is a hard process and the applicability of the perturbation theory is guaranteed by the smallness of distances at which this production takes place ($\sim \frac{1}{2m_Q}$). It is natural to consider that the heavy quarks, from which the heavy-quarkonia bound states are built of, are produced from the collisions of light quarks and gluons from interacting hadrons (we neglect the "sea" content of heavy quark flavors in hadrons). However the main trouble encountered in the quarkonium production remains the problem as to how the produced free heavy quarks join into a bound system. In models of semi-local duality type [1-7], it is assumed that the free $Q, \bar{Q}$ quarks are produced via subprocesses of the type: $q\bar{q} \longrightarrow Q\bar{Q}$ and $gg \longrightarrow Q\bar{Q}$, while the bound state production from Q and $\bar{Q}$ is described as follows: $\sigma(q\bar{q}, gg \rightarrow (Q\bar{Q})) = F \cdot \int_{4m_Q^2}^{4M_Q^2} dQ^2 \cdot \sigma(q\bar{q}, gg \rightarrow Q\bar{Q})$.

Here the integration goes over Q^2 - squared invariant mass of the $Q\bar{Q}$ pair whose upper bound is $2M_Q$, i.e. two masses of a lightest meson carrying the Q flavor (for the c-quark, e.g., it is the mass of $\mathcal{D}$-meson).

One can see that this model has a number of serious drawbacks. So, it does not ensure colorlessness of bound state (e.g., in $q\bar{q}$ annihilation to quarkoinium); the decolorization is assumed to proceed due to the soft gluons emitted by the heavy quarks before the colorless bound state formation. In addition, this model does not fix the spin-parity structure of bound state; e.g., the heavy quarks produced from two gluons cannot be bound into the vector $^3S_1$ quarkonium. All these uncertainties are driven into the factor F which is a free parameter. By fitting parameter F one can obtain the J/Ψ and Υ mesons production cross sections more or less agreeing with experimental data [8].

In order to get rid of all these difficulties, a model was suggested [9-12] in which each of possible bound states of $Q\bar{Q}$ system was determined by a wave function containing all quantum numbers (spin, orbital momentum, charge parity, colorlessness) of the studied bound state. Then, different subprocesses will play different roles in the $^{2S+1}L_J$ quarkonia production. Thus, in the $^3S_1$ and $^1P_1$ quarkonia production the contribution comes from the third order in α_s, $gg \to {}^3S_1({}^1P_1)g$ diagrams (see Fig.1), whereas the second order in α_s, $gg \to {}^1S_0\ ({}^3P_{0,2})$ diagrams are sufficient for the C-even quarkonia production ($^1S_0$, $^3P_{0,2}$). Note, however, that while studying the quarkonia production with large transverse momenta (what is interesting and convenient from the experimental viewpoint) we, with necessity, turn, for the production of quarkonia of $^1S_0$ and $^3P_{0,2}$ types, to the third order in α_s diagrams (see Fig.1) which just make the main contribution to the production cross section of quarkonia with large transverse momenta.

The $^3S_1$ quarkonia production (J/Ψ, Υ, T) was studied earlier in Refs. [10-16], and the C-even quarkonia production, $^1S_0$, $^3P_{0,1,2}$ - in Refs. [14, 16-21]. In this work we have studied the production of $^1P_1$

heavy quarkonia in pp and pp collisions. Note beforehand that the results obtained in this work are the same both for pp and $p\bar{p}$ collisions, since the contribution to the ${}^1P_1$ quarkonia production in the perturbation theory lowest order comes only from subprocesses of the $gg \to {}^1P_1 g$ type see Fig.1), and the content of gluons is the same in protons and antiprotons.

Before turning to the calculations and the description of the obtained results we note that our consideration is based, first, on the QCD perturbation theory; the ${}^1P_1$ quarkonia production is described by the third order in α_s diagrams (see Fig.2), where initial gluons are hadron constituents, while the final gluon yields a hadronic jet. Second, the ${}^1P_1$ - quarkonium is considered as a nonrelativistic bound system of heavy quark and antiquark, and its Bethe-Salpeter wave function (see, e.g., [22,23]) is constructed with account of the laws of addition of spins, angular momenta, and charge parity and colorlessness of the studied bound state.

The expression for the amplitude of the studied process (see Fig.2) can be written as follows [22,23] :

$$A = \int \frac{d^4 q}{(2\pi)^4} \, Sp \left[M(q) \, \chi(P, q) \right] , \tag{1}$$

where P_μ is a 4-momentum of quarkonium, q_μ is a relative 4-momentum of quark and antiquark, $\chi(P, q)$ is the Bethe-Salpeter wave function of ${}^1P_1$ - quarkonium, $M(q)$ is the amplitude corresponding to diagrams of Fig.2.

For the studied process we have:

$$M(q) = g^3 Sp\left(\frac{\lambda^c}{2} \frac{\lambda^b}{2} \frac{\lambda^a}{2}\right) \frac{\gamma^\mu \left[\frac{\hat{P}}{2} + \hat{q} - \hat{K}_1 + m \right] \gamma^\nu \left[-\frac{\hat{P}}{2} + \hat{q} - \hat{K}_3 + m \right] \gamma^\lambda}{\left[\left(\frac{P}{2} + q - K_1\right)^2 - m^2 \right] \left[\left(-\frac{P}{2} + q - K_3\right)^2 - m^2 \right]} \times \tag{2}$$

$$A^a_\mu(K_1)\, A^b_\nu(K_2)\, A^c_\lambda(K_3) + \text{contributions from cyclic permutations} \begin{pmatrix} 1 \to 2 \to 3 \\ a \to b \to c \\ \mu \to \nu \to \lambda \end{pmatrix}$$

Here g is the strong-interaction constant; m is a mass of heavy quark Q ; $a, b, c = 1,2,\dots,8$; λ^a are the Gell-Mann matrices; K_1 , K_2 and K_3 are 4-momenta of gluons, while $A^a_\mu(K_1)$, $A^b_\nu(K_2)$ and $A^c_\lambda(K_3)$ are their polarization 4-vectors. Amplitude $\chi(P,q)$ in the nonrelativistic limit has the form (see, e.g., [22]):

$$\chi(P,q;J,J_z,L,S)=\sum_{M,S_z}(2\pi)\delta\left(q^0-\frac{\vec{q}^{\,2}}{2m}\right)\Psi_{LM}(\vec{q})P_{SS_z}(P,q)\times\langle LM;SS_z|JJ_z\rangle, \tag{3}$$

$$P_{SS_z}(P,q)=\sqrt{\frac{1}{m}}\sum_{s\bar{s}}u\left(\frac{P}{2}+q;s\right)v\left(\frac{P}{2}-q,\bar{s}\right)\langle\tfrac{1}{2}s;\tfrac{1}{2}\bar{s}|SS_z\rangle,$$

where $s, \bar{s}$ are the quark and antiquark spins; S, L, J are respectively spin, orbital momentum and total angular momentum of quarkonium; $\Psi_{LM}(\vec{q})$ is the momentum part of wave function of quarkonium with account of its orbital momentum; $P_{SS_z}(P,q)$ is spin projection operator. The δ function arises because the relative momentum q is small in the nonrelativistic limit, so Q and $\bar{Q}$ quarks can be regarded free and being on the mass shell: $\left(\frac{P}{2}+q\right)^2-m^2\simeq 0$ and in the quarkonium rest frame ($\vec{P}=0$, $P^0=2m$) we obtain $q^0\simeq\frac{\vec{q}^{\,2}}{2m}$.

For the P -wave state of ${}^1P_1$ the operator $P_{SS_z}(P,q)$ is as follows [22] (${}^1P_1$ is a state with total spin 0):

$$P_{00}(P,q)=\sqrt{\frac{1}{8m^3}}\left[-\left(\frac{\hat{P}}{2}+\hat{q}\right)+m\right]\gamma_5\left[\frac{\hat{P}}{2}+\hat{q}+m\right]. \tag{4}$$

Expanding expression (1), with the use of formulae (2) and (4), in powers of small parameter q/m , and extracting the linear terms for the P -wave state of ${}^1P_1$ production we obtain

$$A({}^1P_1) = \sqrt{\frac{3}{4\pi M}}\, R_p'(0)\, Sp\left[-\varepsilon^{\alpha} M^{\alpha}(0)\gamma_5 \frac{\hat{p}+M}{2} + M(0)\gamma_5 \frac{\hat{\varepsilon}\,\hat{p}}{M}\right], \quad (5)$$

where M is the ${}^1P_1$ quarkonium mass which in the nonrelativistic limit is equal to $2m$, ε^{α} is the ${}^1P_1$ quarkonium polarization 4-vector.

$$M^{\alpha}(0) = \left.\frac{\partial M(q)}{\partial q_{\alpha}}\right|_{q=0}, \qquad \int \frac{d^3q}{(2\pi)^3}\, q^{\alpha}\, \Psi_{1M}(\vec{q}) = -i R_p'(0) \sqrt{\frac{3}{4\pi}}\, \varepsilon^{\alpha}.$$

One can see that amplitude A is proportional to the wave function derivative $R_p'(0)$ at zero, this being a result of the fact that the state under study is the P-wave one. A final expression for the amplitude A has the form:

$$A = -8ig^3 \sqrt{\frac{1}{4\pi M}} \left(\frac{1}{4} d_{abc}\right) R_p'(0)\, A_{\mu}^{a}(K_1)\, A_{\nu}^{b}(K_2)\, A_{\lambda}^{c}(K_3)\, \varepsilon_{\alpha} \times$$
$$\times \left\{ \frac{1}{(PK_1)(PK_2)} \left[g_{\mu\alpha} K_{2n} P_m \varepsilon^{\lambda n \nu m} - g_{\nu\alpha} K_{1n} P_m \varepsilon^{\mu n \lambda m} \right] + \right. \quad (6)$$
$$+ \frac{1}{(pK_1)(pK_2)} \left[\frac{K_2^{\alpha}}{(pK_2)} - \frac{K_1^{\alpha}}{(pK_1)} \right] \left[-(pK_2) K_{1\beta} \varepsilon^{\lambda\beta\mu\nu} + g_{\mu\lambda} K_{1n} K_{2\varkappa} P_{\beta} \varepsilon^{n\varkappa\nu\beta} - \right.$$
$$- K_{1\lambda} K_{2n} P_{\beta} \varepsilon^{\mu n \nu\beta} - K_{2n} P_{\mu} P_{\beta} \varepsilon^{\beta\lambda n\nu} + K_{1n} K_{3\varkappa} P_{\nu} \varepsilon^{\lambda n \mu\varkappa} \Big] +$$

$+$ terms from permutations $\begin{bmatrix} 1 \to 3 \\ \mu \to \lambda \end{bmatrix}$ and $\begin{bmatrix} 2 \to 3 \\ \nu \to \lambda \end{bmatrix}$.

Then for the cross section we obtain:

$$E\frac{d\sigma}{d^3p} = \frac{z_0}{\pi}\frac{d\sigma}{dp_\perp^2 dz_0} = \frac{5}{3\cdot 2^9 M}\int dx_1 \frac{\alpha_s^3}{x_1 x_2 S_0^2}\frac{G(x_1)G(x_2)}{(x_1 - z_0)} A_0^2 |R_P'(0)|^2, \quad (7)$$

here E is the quarkonium energy, and $p_\perp^2$ is its squared transverse momentum. $z_0 = \frac{(PP_2)}{(P_1P_2)}$, where P_1 and P_2 are momenta of initial hadronic beams, x_1 and x_2 are momenta fractions carried away by gluons from hadrons ($K_1 = x_1 P_1$) and $K_2 = x_2 P_2$), $\alpha_s = \frac{g^2}{4\pi}$, $S_0 = (P_1 + P_2)^2$, $G(x)$ are distribution functions of gluons in hadrons, A_0^2 is the squared expression in curly brackets of formula (6). Since we are interested not only in the total cross section of $^1P_1$ quarkonium production but also in its differential distributions over $p_\perp^2$ and z_0, we'll write out expressions for the invariant four-dimensional products which enter A_0^2 through $p_\perp^2$, z_0 and S_0 :

$$\begin{aligned}
&(K_1P) = \frac{1}{2}\left[M^2 + (x_1 - z_0)x_2 S_0\right], \quad (K_2P) = \frac{1}{2}x_2 z_0 S_0, \\
&(K_3P) = \frac{1}{2}\left[x_1 x_2 S_0 - M^2\right], \quad (K_1K_2) = \frac{1}{2}x_1 x_2 S_0, \\
&(K_1K_3) = \frac{1}{2}\left[x_2 z_0 S_0 - M^2\right], \quad (K_2K_3) = \frac{1}{2}(x_1 - z_0)x_2 S_0, \qquad (8) \\
&P^2 = M^2, \quad K_1^2 = K_2^2 = K_3^2 = 0, \\
&x_2 = \frac{x_1 p_\perp^2}{z_0(x_1 - z_0)S_0} + \frac{M^2}{z_0 S_\perp},
\end{aligned}$$

as well as the variation intervals of the variables x_1, $p_\perp^2$ and z_0 :

$$\begin{aligned}
&\frac{z_0(z_0S_0 - M^2)}{z_0 S_0 - p_\perp^2 - M^2} \leqslant x_1 \leqslant 1, \\
&0 \leqslant p_\perp^2 \leqslant (1 - z_0)(z_0S_0 - M^2), \qquad (9) \\
&\frac{m^2}{S_0} \leqslant z_0 \leqslant 1 .
\end{aligned}$$

Note that in problems with the production of P -wave quarkonia and their decays to gluons, there arises a problem with infrared divergences, since the soft gluon which is responsible for that always carries away the orbital momentum [24]. To get rid of this divergence, it is enough to introduce a cut-off at lower bound of $p_\perp^2$ integration: $p_\perp^2 \geqslant p_{\perp\,min}^2 \simeq$ $\simeq$ a few GeV^2. This will provide the difference of z_0 from 1 (in fact, in the lab system, when one of initial hadrons is in rest, $z_0 = E^L/E_1^L$, i.e. z_0 is a portion of energy carried away by quarkonium from the hadron initial beam).

To obtain the total and differential cross sections from formulae (7-9), one should have an information on α_s , $|R'_P(0)|^2$, M of the studied $^1P_1$ quarkonia. For masses of charmonium and bottomonium we take an average value of masses of the corresponding χ -even quarkonia: $M_{c\bar c}(^1P_1) \simeq 3.5$ GeV, $M_{b\bar b}(^1P_1) \simeq 9.9$ GeV. For toponium mass we take two values: $M_{1t\bar t}(^1P_1) = 45.5$ GeV and $M_{2t\bar t}(^1P_1) = 80$ GeV. For α_s we take the following values: $\alpha_s(M^2_{c\bar c}) = 0.4$, this being in agreement with experimental data [25]; then, according to the logarithmic scale of α_s variation, we have: $\alpha_s(M^2_{b\bar b}) = 0.26$, $\alpha_s(M^2_{1t\bar t}) = 0.19$ and $\alpha_s(M^2_{2t\bar t}) = 0.17$.

Introduce a notation

$$\tau = \frac{4\,|R'_P(0)|^2}{|R_S(0)|^2\,(M(^3S_1))^2}\,, \tag{10}$$

while the value of $R_S(0)$ we derive from the lepton widths of their decays:

$$\Gamma_{ee} = \frac{4\alpha^2 e_Q^2}{(M(^3S_1))^2}\,|R_S(0)|^2\,,$$

where e_Q is the heavy quark charge; α is the fine structure constant; Γ_{ee} is the width of vector quarkonium decay to e^+e^- pair (Γ_{ee} $\cong$

$\equiv \Gamma\,({}^3S_1 \to e^+e^-))$. The ratio τ depends on the choice of potential model; so, in the models of the [26-28] type there is predicted: $\tau = 0.074$ for charmonium, $\tau = 0.013$ for bottomonium and $\tau = 0.001$ for toponium. Besides, the ratio $\Gamma_{ee}/\mathcal{L}_Q^2$ in these models is a constant mass-independent quantity. Then Γ_{ee} for the toponium decay is assumed 4.8 keV (just as for charmonium).

Note beforehand that the results obtained by us for toponium depend essentially on the model choice. Indeed, the assumption that $\Gamma_{ee}/\mathcal{L}_Q^2$ is mass-independent [26-28] (which works well enough for charmonium and bottomonium), when kept to the toponium masses may work poorly since in the toponium system the Coulomb interaction may play a large role, and for the Coulomb potential $\Gamma_{ee} \sim M$ and then $\Gamma_{ee}(T) \gg \Gamma_{ee}(J/\Psi)$. This undoubtedly will increase our obtained ${}^1P_1$ toponium production cross sections at least by an order of magnitude. In the Richardson model (see [29]), $\Gamma_{ee}(T)$ is only twice as large as $\Gamma_{ee}(J/\Psi)$; however in this model τ exceeds the value used by us for toponium more than by an order of magnitude. I.e. one can see that the predicted cross section of the ${}^1P_1$ toponium production depends greatly on the choice of the model.

To estimate the cross sections, we must also know the gluon distribution function. For that we take the standard scaling function $x \cdot G(x) = 3 \cdot (1-x)^5$. Besides, in order to check the sensitivity of obtained results to $G(x)$ we shall use a nonscaling function:

$$x\,G(x) = (2{,}01 - 2{,}73\rho + 1{,}29\rho^2)\,x^{-0{,}93\rho + 0{,}36\rho^2}(1-x)^{2{,}3+1{,}83\rho}, \qquad (11)$$

where $\rho = \ln\left[\dfrac{\ln(M^2+P_\perp^2)/\Lambda^2}{\ln Q_0^2/\Lambda^2}\right]$, $\Lambda = 0.2$ GeV, $Q_0^2 = 5$ GeV^2, taken from Ref. [30] ; here the strong-interaction constant α_s is taken in the form:

$$\alpha_s = \frac{12\pi}{(33-2n)\ln\left((M^2+P_\perp^2)/\Lambda^2\right)} , \qquad (12)$$

where n is the number of included flavors. Note that the use of another scaling distribution function, $xG(x) = (1-x)^5\,(1.07 + 13.50x)$, obtained from the analysis of the deep inelastic leptonic processes in Ref.[31], does not bring to a considerable change of results obtained by us (the change reaches $10 \div 20$ %).

Figs. 3 and 4 show the total cross section of ${}^1P_1$ quarkonium production versus $\sqrt{s_0}$. Fig. 3 presents the total cross section of ${}^1P_1$ -charmonium and bottomonium production, while in Fig. 4 the cross section of ${}^1P_1$ - toponium production for two masses $M = 45.5$ and 80 GeV is given. One can see that the cross sections of charmonium and bottomonium production come out to the asymptotic regime at $\sqrt{s_0} \sim 500$ GeV, while that of toponium - at $\sqrt{s_0} \sim 3$ TeV. All the curves are given for the standard gluon distribution function. The use of the nonscaling distribution function according to formula (11) results in a rather great change of obtained cross sections: they increase twice as much. It is interesting also to note that the cross section of ${}^1P_1$ charmonium and bottomonium production is by two orders of magnitude as less as the cross section of corresponding ${}^3S_1$ quarkonia production [16] (this comparison is interesting, since similar diagrams contribute to the production of ${}^3S_1$ and ${}^1P_1$ quarkonia). This undoubtedly is due to the smallness of τ (in our used models $\tau_c = 0.074$ and $\tau_b = 0.013$). For toponium this difference reaches already three orders ($\tau_t = 0.001$). Emphasize once again that the value of the cross section of ${}^3S_1$ and ${}^1P_1$ - toponia production depends essentially on the model choice and may be larger at least by an order of magnitude for ${}^1P_1$ -toponium. One can see from the obtained results that since the dynamics of $^3S_1$ and $^1P_1$ -quarkonia is the same, and the difference in the production cross sections almost entirely is determined by the fact that $^3S_1$ is the S -wave, and $^1P_1$ is the P -wave bound states, then the role of the fact that $^3S_1$ is the vector, and $^1P_1$ is the pseudovector quarkonia has turned out insignificant. I.e. the dynamics of the studied process (see Fig. 2) has completely suppressed the effect of difference of vertices of quark and antiquark transition to the bound system (see formula (5) of this paper and (18) of Ref. [22]).

In Figs. 5-8 the differential characteristics for the $^1P_1$ -quarkonium production cross section are given. Fig. 5 shows the dependence of $^1P_1$ -charmonium production cross section, $E\frac{d\sigma}{d^3p}$, on squared transverse momentum of produced quarkonium $p_\perp^2$ at different values of z_0 (for $\sqrt{S_0}$ = 63 and 540 GeV). Fig. 6 presents the analogous cross section for $^1P_1$ -bottomonium (for $\sqrt{S_0}$ = 63, 540 and 2000 GeV), and in Figs. 7 and 8 the same is given for $^1P_1$ -toponium (where $\sqrt{S_0}$ = 540 and 2000 GeV) at two values of masses: M = 45.5 and 80 GeV. One can see from the figures that the cross sections decrease rather rapidly with increasing transverse momentum, the rate of decrease corresponding to what took place for $^3S_1$ -quarkonia [14]. This undoubtedly is connected with the same production mechanisms of $^3S_1$ and $^1P_1$ quarkonia. This result differs from what took place when comparing the $^1S_1$ production, on the one hand, and $^1S_0$ and $^3P_{0,1,2}$ -quarkonia production, on the other [14]. The spectra of $^1S_0$ and $^3P_{0,1,2}$ quarkonia production were flattening [14,19], i.e. they decreased more slowly with increasing $p_\perp$, this being due to the fact that not only diagrams of Fig.2 contribute to the production of these states. Note also some slowing-down in the rate of decrease of differential cross section of $^1P_1$ -toponium production with increasing $p_\perp^2$ *relative to* $^1P_1$ *-charmonium and bottomonium (see Figs. 5-8).* As one can see, all the cross sections decrease greatly with increasing z_0 .

We have studied the production of ${}^1P_1$ -quarkonia (charmonium, bottomonium and toponium) in the hard processes of pp- and $p\bar{p}$-collisions for energies from ISR ($\sqrt{S_0} = 63$ GeV) and CERN ($\sqrt{S_0} = 540$ GeV) up to the TEVATRON ($\sqrt{S_0} = 2$ TeV) and UNK ($\sqrt{S_0} = 4 + 6$ TeV) energies. Cross sections for these quarkonia production are obtained. Thus, the ${}^1P_1$ -charmonium is produced with a cross section $\sigma({}^1P_1(c\bar{c})) \simeq 4 \cdot 10^4$ pb at $\sqrt{S_0} \simeq 1$ TeV, while $\sigma({}^1P_1(b\bar{b})) \simeq 3 \cdot 10^2$ pb and $\sigma({}^1P_1(t\bar{t})) \simeq 10^{-1}$ pb at the same energies. At luminosities ($\mathcal{L} \simeq 10^{31}$ cm^{-2}s^{-1}) being planned on TEVATRON and UNK we have a sufficiently rich production of ${}^1P_1$ -charmonium and bottomonium. The detection of these states may proceed over ${}^1P_1 \to J/\Psi(\Upsilon) + \pi^+ + \pi^-$ or ${}^1P_1 \to J/\Psi(\Upsilon) + \pi^0 + \pi$ decays which have 1% branching, as well as over ${}^1P_1 \to {}^1S_0 + \gamma$ proceeding with large branching (50%) with a subsequent ${}^1S_0 \to 2\gamma$ decay. Note, in addition, that in the studied mechanism the ${}^1P_1$ -quarkonia production proceeds in association with hard hadronic jet, which also will provide the identification of ${}^1P_1$ -states. In case of ${}^1P_1$ -toponium production the cross section is, of course, very small. Even if it is by an order higher (we have noted that it strongly depends on the choice of the model), at planned luminosities we shall have for ${}^1P_1$ -toponium production 100 events per year.

The present analysis of ${}^1P_1$ -quarkonia production was stimulated by the fact that these states are not found in the cascade decays $\Psi' \to \pi^0\,{}^1P_1 \to \pi^0 \eta_c \gamma$ (see, e.g., 32). The search for ${}^1P_1$ state in $\Psi' \to \chi_2\gamma \to \gamma\gamma\,{}^1P_1$ and $\Psi' \to \eta_c'\gamma \to \gamma\gamma\,{}^1P_1$ cascades is not carried out because of the smallness of expected width of these decays [33]. The hadronic transition $\Upsilon(3S) \to {}^1P_1 + \pi + \pi \to {}^1S_0 + \gamma + \pi + \pi$ was expected to be promising. According to the estimation in [34], the probability for this cascade amounts to one percent. However, in Ref. [35] it is shown that the obtained estimate for the probability of this decay, resulting from the parton model (where it is considered that $\Upsilon(3S) \to {}^1P_1 + 2g$ and then the real gluons transform to π-mesons) is incorrect, since the energy corresponding to gluons is low (~ 400 MeV), so the gluons cannot be considered as hard. The account of soft gluon effect [35] has reduced the probability of this transition by two orders of magnitude. Thus we see that the observation of ${}^1P_1$-quarkonia production is problematic in various transitions. We'd like to point out also the attempt to detect the ${}^1P_1$-charmonium at the ISR [36] in direct production of charmonium in $p\bar{p}$-annihilation. The ${}^1P_1$-quarkonium identification was carried out by ${}^1P_1 \to J/\Psi + \pi + \pi$ decay, where several events were interpreted as ${}^1P_1$-charmonium production. If however taking into consideration that according to [35] this cascade is suppressed, then it is quite unclear where such an intense signal with $J/\Psi + \pi + \pi$ comes from.

It is worth mentioning one more possibility of ${}^1P_1$-quarkonium production in $p\bar{p}$-, pp-collisions, when the 2^1S_0-quarkonium is produced [14] (see Fig.1) with a subsequent decay with ${}^1P_1$ production ($2^1S_0 \to {}^1P_1 + \gamma$). The ${}^1P_1$ quarkonium production via 2^1S_0-quarkonium is significant, since the cross section of the 2^1S_0-quarkonium production exceeds by three orders the cross section of the ${}^1P_1$-quarkonium direct production, and despite the smallness of $B(2^1S_0 \to {}^1P_1 + \gamma) \sim 10^{-3}$ [37] due to small difference between masses of 2^1S_0 and ${}^1P_1$ states, the ${}^1P_1$ quarkonia production is of the same order and even larger than its direct production. It is interesting that since the cross section of 1^1S_0 and 2^1S_0 quarkonia decreases with increasing $p_\perp$ more slowly than the cross section of ${}^3S_1$ and ${}^1P_1$ quarkonia [14,12], the production of ${}^1P_1$ quarkonium from the decay of the 2^1S_0 will be significant rather for very large transverse momenta than from the direct production (Fig.2).

In conclusion, we'd like to discuss a possibility of ${}^1P_1$ quarkonium production in photon-hadron collisions. Then the main contribution to ${}^3S_1$ and ${}^1P_1$ quarkonia production will come from diagrams of Fig.2, where one of the initial gluons is to be replaced by the photon. The study of heavy quarkonia production in γp collisions is interesting just for the reason that the ${}^1S_0$ and ${}^3P_{0,1,2}$-quarkonia production here cannot proceed in the lowest order of the perturbation theory (see Fig.1, where one of the initial gluons is replaced by the photon). One can readily see that the colored factor cancels the contributions from all diagrams, while diagrams of the Fig.2 type do not contribute to the ${}^1S_0$ and ${}^3P_{0,1,2}$ because of the C-parity, this being different from the situation we have with the hadron-hadron collisions. However, for the ${}^1P_1$ quarkonium photon production we have obtained cross sections which are at least by two orders lower than those we have in the pp and $p\bar{p}$ collisions. This undoubtedly makes problematic the search even for ${}^1P_1$ -charmonium and bottomonium in these reactions.

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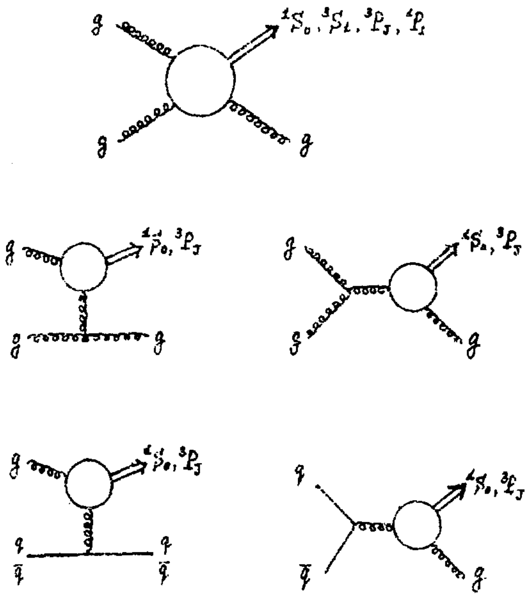
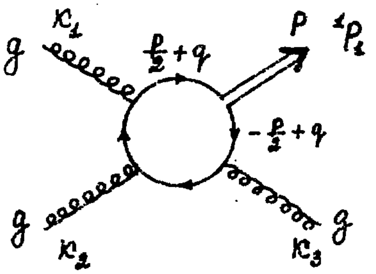


Fig. 1



+ diagrams with permutations

Fig. 2

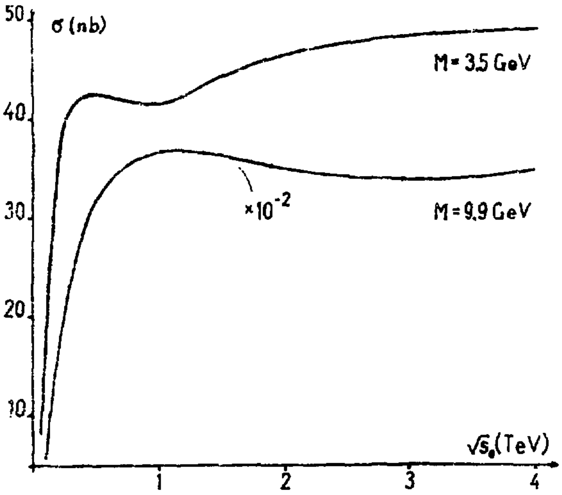

σ (nb)
50
40
30
20
10
M = 3.5 GeV
$\times 10^{-2}$
M = 9.9 GeV
$\sqrt{s_0}$ (TeV)
1
2
3
4

Fig. 3

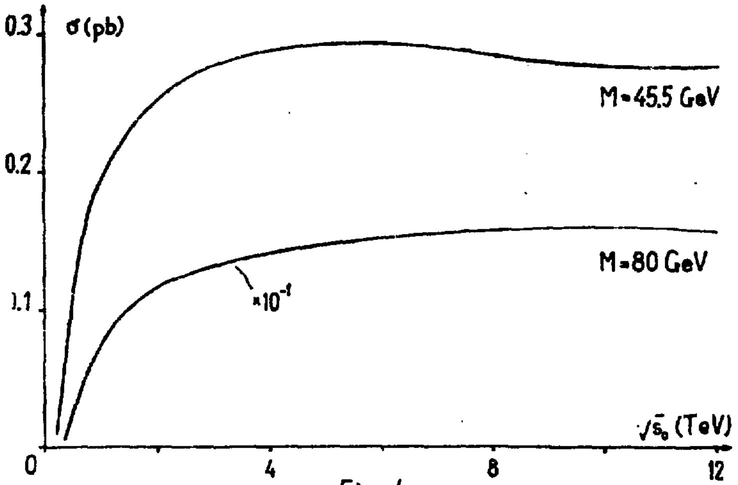

0.3
σ (pb)
0.2
0.1
M = 45.5 GeV
M = 80 GeV
$\times 10^{-1}$
$\sqrt{s_0}$ (TeV)
0
4
8
12

Fig. 4

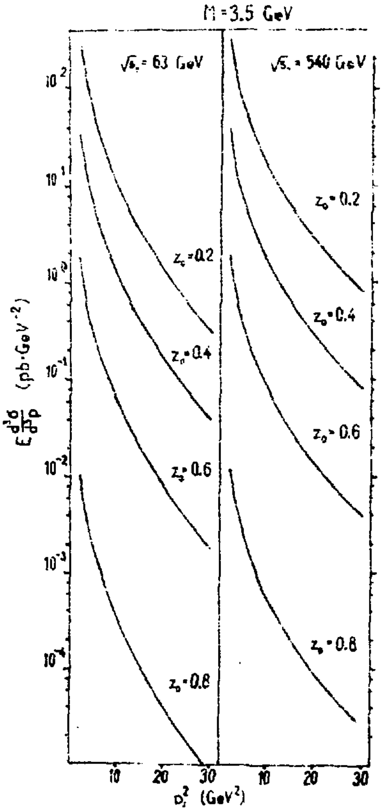

M = 3.5 GeV
√s₀ = 63 GeV
√s₀ = 540 GeV
10²
10¹
10⁰
10⁻¹
10⁻²
10⁻³
10⁻⁴
E d³σ/d³p (pb·GeV⁻²)
z₀ = 0.2
z₀ = 0.4
z₀ = 0.6
z₀ = 0.8
z₀ = 0.2
z₀ = 0.4
z₀ = 0.6
z₀ = 0.8
10
20
30
10
20
30
p⊥² (GeV²)

Fig. 5

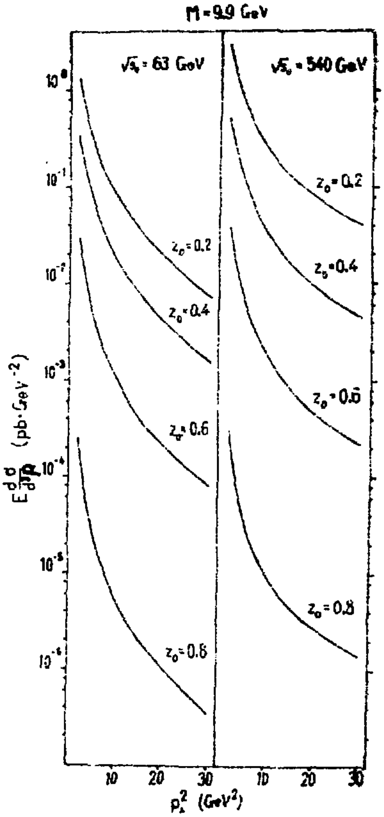

M = 9.9 GeV
√s₀ = 63 GeV
√s₀ = 540 GeV
10⁰
10⁻¹
10⁻²
10⁻³
10⁻⁴
10⁻⁵
10⁻⁶
E dσ/d³p (pb·GeV⁻²)
z₀ = 0.2
z₀ = 0.4
z₀ = 0.6
z₀ = 0.8
z₀ = 0.2
z₀ = 0.4
z₀ = 0.6
z₀ = 0.8
10
20
30
10
20
30
p⊥² (GeV²)

Fig. 6

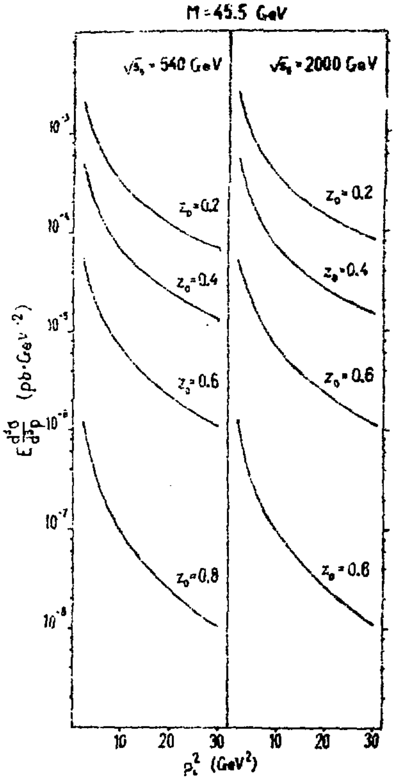

M=45.5 GeV
$\sqrt{s_0}=540$ GeV
$\sqrt{s_0}=2000$ GeV
10^{-3}
10^{-4}
10^{-5}
10^{-6}
10^{-7}
10^{-8}
$E\frac{d^3\sigma}{d^3p}$ (pb·GeV^{-2})
$z_0=0.2$
$z_0=0.4$
$z_0=0.6$
$z_0=0.8$
$z_0=0.2$
$z_0=0.4$
$z_0=0.6$
$z_0=0.8$
10
20
30
10
20
30
p_t^2 (GeV^2)

Fig. 7

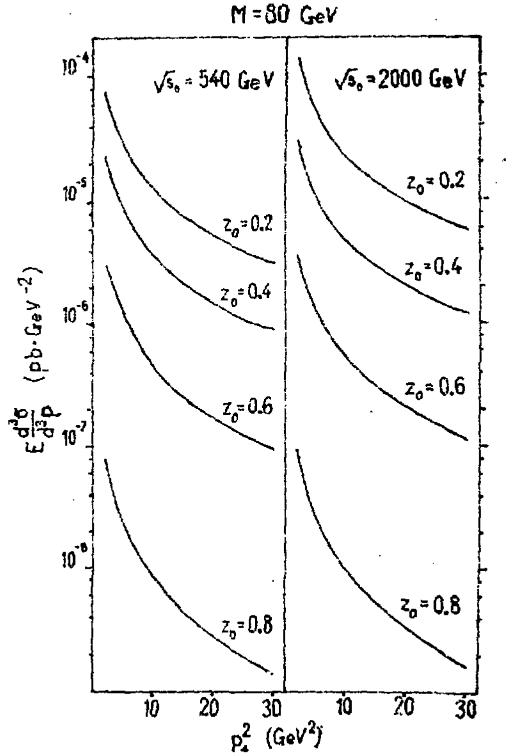

M=80 GeV
$\sqrt{s_0}=540$ GeV
$\sqrt{s_0}=2000$ GeV
10^{-4}
10^{-5}
10^{-6}
10^{-7}
10^{-8}
$E\frac{d^3\sigma}{d^3p}$ (pb·GeV^{-2})
$z_0=0.2$
$z_0=0.4$
$z_0=0.6$
$z_0=0.8$
$z_0=0.2$
$z_0=0.4$
$z_0=0.6$
$z_0=0.8$
10
20
30
10
20
30
p_t^2 (GeV^2)

Fig. 8

Figure Captions

Fig. 1. **Diagrams that contribute to the heavy-quarkonia production at a quark-gluon level.**

Fig. 2. Diagrams contributing to the $^1P_1$ -quarkonia production.

Fig. 3. The total cross section of the $^1P_1$ -charmonium and bottomonium production versus $\sqrt{S_o}$.

Fig. 4. The total cross section of the $^1P_1$ -toponium production versus $\sqrt{S_o}$, given for two masses: M = 45.5 and 80 GeV.

Fig. 5. The dependence of twice-differential cross section $E\frac{d\sigma}{d^3p}$ of the $^1P_1$ charmonium production on $p_\perp^2$ for $\sqrt{S_o}$ = 63 and 540 GeV and different values of z_o .

Fig. 6. The $E\frac{d\sigma}{d^3p}$ cross section of the $^1P_1$ -bottomonium production as a function of $p_\perp^2$ for $\sqrt{S_o}$ = 63 and 540 GeV and different values of z_o .

Fig. 7. The $E\frac{d\sigma}{d^3p}$ cross section of the $^1P_1$ -toponium with M = 45.5 GeV as a function of $p_\perp^2$ for $\sqrt{S_o}$ = 540 and 2000 GeV and different values of z_o .

Fig. 8. The $E\frac{d\sigma}{d^3p}$ cross section of the $^1P_1$ -toponium production with M = 80 GeV as a function of $p_\perp^2$ for $\sqrt{S_o}$ = 540 and 2000 GeV and different values of z_o .

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РОЖДЕНИЕ $^1P_1$ ТЯЖЕЛЫХ КВАРКОНИЕВ В АДРОН-АДРОННЫХ СТОЛКНОВЕНИЯХ

(на английском языке, перевод Асланян З.Н.)

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ЕРЕВАНСКИЙ ФИЗИЧЕСКИЙ ИНСТИТУТ