ИНСТИТУТ ЯДЕРНОЙ ФИЗИКИ СО АН СССР

V.S. Fadin and O.I. Yakovlev

NONPERTURBATIVE CORRECTION TO THE THRESHOLD PRODUCTION OF $t\bar{t}$-PAIRS

PREPRINT 90-138

НОВОСИБИРСК

Nonperturbative Correction to the Threshold Production of $t\bar{t}$-Pairs

V.S. Fadin and O.I. Yakovlev

Institute of Nuclear Physics
630090, Novosibirsk, USSR

ABSTRACT

Calculations are performed of the nonperturbative correction to the cross-section for the $t\bar{t}$-pair creation near the threshold in e^+e^--annihilation, which is connected with the existence of a gluon condensate. These have been made using a constant chromoelectric field approximation.

1. Despite of the t-quark is not yet discovered experimentally, its existance is beyond doubt. So far rather strong limitations on the mass of this quark have been obtained: $m_t = 137 \pm 40$ GeV [1]. These limitations are derived from the whole of high energy physics experiments: a direct search for the t-quark, observations of large $B_d^0 - \bar{B}_d^0$ mixing, measurement of the Z^0-boson mass m_Z and its hadron and lepton widths on the SLC and the LEP collider, of the W-boson mass m_W and the ratio m_W/m_Z on hadron colliders, as well as from neutrino scattering data. The large value of m_t leads to the fact that the physical properties of the t-quark are drastically different from those of its predecessors — the c- and b-quarks. First of all, this is due to a large width of the t-quark. For $m_t > m_W$ the semiweak decay $t \to W^+ + b$ channel opens. The widths of this decay rapidly increases with increasing m_t. At $m_t = 100$ GeV it equals about 100 MeV, and at $m_t = 110$ GeV it becomes roughly equal to the width of the ρ-meson and for m_t above 150 GeV it reaches an asymptotical regime [2]:

$$\Gamma_t = (180 \text{ MeV}) \cdot \left(\frac{m_t}{m_W}\right)^3. \qquad (1)$$

The large width of the t-quark radically changes the cross-section for $e^+e^- \to t\bar{t}$ process near the threshold. The near-threshold region of heavy quarks production in e^+e^--annihilation offers ample scope for studying the properties of quarks and it is usually used in analysis. The cross-section for the $t\bar{t}$-pair creation with an allowance for the quark width has been considered in detail in terms of the perturbative QCD in Refs [3—5]. The goal of the present paper is to calculate the nonperturbative correction to the cross-section for $e^+e^-\to t\bar{t}$ process.

2. The correction to the cross-section we are going to calculate arises because of the existence of long-wave fluctuation of the gluon field in QCD vacuum [6]. The characteristic space-time sizes of these fluctuations are usually estimated as $r_f\sim\tau_f\sim\Lambda_{QCD}$, where $\Lambda_{QCD}\sim 100$ MeV is the dimensional parameter of QCD.

The method of calculating the nonperturbative effects in a nonrelativistic quarkonium was suggested in [7, 8]. It is based on a multipole expansion of the interaction between the quarkonium and the gluon field of nonperturbative fluctuations. Indeed, if $r_{t\bar{t}}$ and $\tau_{t\bar{t}}$ are the characteristic sizes and the times of evolution of the quark system and the inequalities

$$\frac{r_{t\bar{t}}}{r_f}\ll 1, \tag{2}$$

$$\frac{\tau_{t\bar{t}}}{\tau_f}\ll 1 \tag{3}$$

are satisfied, we can confine ourselves, in the leading order of these parameters, to the interaction between the colour dipole moment of the $t\bar{t}$-system and the constant chromoelectric field.

In Refs [7, 8], the corrections to the energy levels and to the wave functions of the bound states were calculated in the constant chromoelectric field approximation. In this approximation, the corrections are expressed via the gluon condensate [6, 9]:

$$\left\langle \frac{\alpha_s}{\pi}(G^a_{\mu\nu})^2\right\rangle = 0.012-0.018\ (\mathrm{GeV})^4, \tag{4}$$

where $G^a_{\mu\nu}$ is the gluon field tensor.

We have to note that by virtue of the nonrelativistic nature of the quark system, the characteristic times of its evolution are large as compared with the characteristic distances. For this reason, the condition (3) is restrictive to a more extent than (2). To eliminate it (3), we have to take the time dependence of the chromoelectric field into account. This problem was treated in Ref [10]. It has been solved at the expence of the introduction of nonlocal vacuum condensates [10] the values of which are strongly dependent on the vacuum model.

The corrections, calculated in [7, 8], are applicable only for rather low levels. As the level number grows, the characteristic sizes and the evolution times of the quark system increase and, as a consequence, the conditions (2) and (3) violate. The increase of the characteristic sizes leads to a drastic growth of corrections [7, 8] with the increase of the level number. The conditions (2) and (3), are, therefore, expressed in the limitation on the magnitude of a correction, i. e. in the requirement for its smallness. However this limitation is independent and, generally speaking, not equivalent neither to (2) nor (3). The relation between them is determined by the structure of the gluon condensate and cannot be strictly defined for the time being. Using the nonrelativistic energy $E=\sqrt{S}-2m_t$, we may say that Refs [7, 8] are concerned with the region of negative energies which are far from zero. As for positive energies, i. e. the region above the threshold of production of a quark-antiquark pair, since in the absence of the width the quark would become rather distant it is clear that the nonperturbative interaction is very important for the formation of final states. And only in the case of «smeared» cross-section (see Ref. [11]), we can hope that the nonperturbative effects may be considered as corrections. Under a natural condition that the interval of «smearing» with respect to energy E is less than or about the magnitude of this energy and with the simplest estimations of the characteristic distances and times, conditions (2) and (3) are represented as $\Lambda_{QCD}/\sqrt{m_t|E|}\ll 1$ and $\Lambda_{QCD}/|E|\ll 1$, respectively. They obviously exclude the near-threshold energies.

3. The large width of the t-quark radically changes the situation. At a rather large width, the contribution of a nonperturbative interaction becomes a calculable correction at any energy. This is the case owing to the following circumstances. First, the existence of the width limits both the sizes of the $t\bar{t}$-system produced and the time of its evolution. In this case the applicability conditions (2) and (3) for a multipole expansion take the form:

$$\frac{\Lambda_{QCD}}{(m_t^2(E^2+\Gamma_t^2))^{1/4}}\ll 1 \tag{5}$$

and

$$\frac{\Lambda_{QCD}}{(E^2+\Gamma_t^2)^{1/2}}\ll 1, \tag{6}$$

respectively. As has already been mentioned, the condition (6) is the most limiting one. However, it is also satisfied at any E at $\Gamma_t \gg \Lambda_{QCD}$, thereby allowing the use of the multipole expansion for such Γ_t to calculate the $t\bar{t}$-production cross-section in the whole region near the threshold.

Second, the contribution of nonperturbative effects to the cross-section, calculated in terms of this expansion, turns out to be small (at fairly large Γ_t) also for any E and, as a consequence, it may be considered as the correction. An approximate evaluation of the influence of the gluon condensate gives rise to the following condition of the smallness of nonperturbative effects [5]:

$$\left\langle \frac{\alpha_s}{\pi}(G^a_{\mu\nu})^2 \right\rangle \ll (E^2+\Gamma_t^2)^{3/2} m_t . \quad (7)$$

The above condition is fullfield for $m_t > 100$ GeV at any E.

4. The $e^+e^- \to t\bar{t}$ cross-section in the near-threshold region is expressed via the imaginary part of the vacuum polarization $\mathscr{P}^{(t)}(s)$ by the vector current of the t-quarks $J^{\mu}_{(t)} = \bar{t}\gamma^{\mu}t$:

$$\sigma_{e^+e^-\to t\bar{t}}(s) = \frac{4\pi\alpha^2}{3s}\, 12\pi \left\{ Q_t^2 + \frac{v_t^2}{\varkappa(1-M_Z^2/s)^2} \right\} \operatorname{Im} \mathscr{P}^{(t)}(s) . \quad (8)$$

Here the first term in the braces corresponds to the contribution of a photon, while the second to that of a Z^0-boson. We neglect the axial part of the Z^0-boson contribution because it contains an additional square of the velocity of the t-quark in comparison with the vector part. In (8), $\sqrt{S} = 2E_e$ is the total energy of colliding beams, $Q_t = 2/3$ — the electric charge of the t-quark and $v_t = 1 - \frac{8}{3}\sin^2\theta_W$ is its vector coupling constant; $\varkappa = (2\sin 2\theta_W)^4$, $\varkappa \simeq 8$ at $\sin^2\theta_W = 0.23$. The terms, proportional to the vector coupling constant of an electron v_e ($v_e \simeq -0.08$) are omitted in (8).

It is well known (see, e. g., [12, 13]) that in the near-threshold region the vacuum polarization $\mathscr{P}^{(t)}(s)$ is expressed via the nonrelativistic Green function of a relative motion of a pair of produced particles. With taking the colour into account,

$$\mathscr{P}^{(t)}(s) = \frac{3}{2m_t^2}\, G_E(0,0) , \quad (9)$$

where $G_E(\vec{r}, \vec{r}\,')$ is the Green function of the $t\bar{t}$-system in a colourless state:

$$G_E(\vec{r}, \vec{r}\,') = \langle \vec{r} \,|\, \hat{G}(E) \,|\, \vec{r}\,' \rangle . \quad (10)$$

Here $\hat{G}(E)$ is the operator which acts in the coordinate space of the $t\bar{t}$-system. For quarks interacting with the gluon field of vacuum fluctuations [7], it is a projection of a more general operation $(\hat{H}-E)^{-1}$ onto the gluon vacuum and onto the colourless $t\bar{t}$-state:

$$\hat{G}(E) = P_s \langle (\hat{H}-E)^{-1} \rangle P_s . \quad (11)$$

The sign $\langle ... \rangle$ in (11) implies an averaging over the gluon vacuum, P_s is the operator of projection onto the colourless $t\bar{t}$-state and $\hat{H}$ is the Hamiltonian of the $t\bar{t}$-system plus the gluon field:

$$\hat{H} = \hat{H}_{t\bar{t}} + \hat{H}_g + \hat{H}_{int} . \quad (12)$$

Here

$$\hat{H}_{t\bar{t}} = P_s \hat{H}_s + P_8 \hat{H}_8 , \quad (13)$$

where

$$\hat{H}_s = \frac{\hat{P}^2}{m_t} - \frac{4}{3}\frac{\alpha_s(r)}{r}; \quad \hat{H}_8 = \frac{\hat{P}^2}{m_t} + \frac{1}{6}\frac{\alpha_s(r)}{r} , \quad (14)$$

$P_8 = 1 - P_s$ is the operator of projection onto the octet state of the $t\bar{t}$-system, $\hat{H}_g$ is the Hamiltonian of the gluon field $\hat{H}_{int}$ is the Hamiltonian of interaction between the quarks and the gluon field of vacuum fluctuations; in a dipole approximation,

$$\hat{H}_{int} = -\frac{g}{2}(t^a - \bar{t}^a)\, \vec{\mathscr{E}}^a(0)\, \vec{r} , \quad (15)$$

where $\mathscr{E}^a_i = G^a_{i0}$, $t^a(\bar{t}^a)$ are the generators of the colour group for a quark (antiquark).

The first nonvanishing correction to the Green function $G_E(\vec{r}, \vec{r}\,')$ caused by the interaction with the vacuum fluctuation field is represented as follows [7]:

$$\Delta G_E(\vec{r}, \vec{r}\,') = -\eta \langle \vec{r} \,|\, \frac{1}{(\hat{H}_s - E)}\, \vec{r}\, \frac{1}{(\hat{H}_8 - E)}\, \vec{r}\, \frac{1}{(\hat{H}_s - E)} \,|\, \vec{r}\,' \rangle , \quad (16)$$

where

$$\eta = \left\langle \frac{\pi\alpha_s}{18} (G^a_{\mu\nu})^2 \right\rangle . \quad (17)$$

At a fixed constant α_s, we have [7]

$$\Delta G_E(0,0) = -\frac{9\eta m_t^3}{2\pi k^5}\lambda^2 \sum_{n=0}^{\infty} \frac{(n+3)!}{n!(n+2+\lambda/8)} \left[\frac{\Gamma(n-\lambda)}{\Gamma(n+5-\lambda)}\right]^2. \tag{18}$$

Here

$$k=\sqrt{-m_t(E+i0)}\ ; \quad \lambda=\frac{2m_t\alpha_s}{3k},$$

$\Gamma(x)$ is the gamma-function.

Formula (9) concerns the case of stable quarks. As has been shown in Refs (3, 4), taking into account of the width of the t-quark reduces to the replacement $E \to E + i\Gamma_t$. With the expression for the cross-section (8) incorporated, we obtain that the nonperturbative correction to the $e^+e^- \to t\bar{t}$ cross-section near the threshold is equal to

$$\Delta\sigma_{e^+e^-\to t\bar{t}}(s) = \frac{4\pi\alpha^2}{3s}\frac{18\pi}{m_t^2}\left[Q_t^2 + \frac{v_t^2}{\varkappa(1-m_Z^2/s)^2}\right\} \operatorname{Im}\Delta G_{E+i\Gamma_t}(0,0), \tag{19}$$

where $\Delta G_{E+i\Gamma_t}(0,0)$ is given by formula (18) with $k = \sqrt{-m_t(E+i\Gamma_t)}$ for the case of fixed α_s.

The cross-sections for e^+e^--annihilation, having the resonance nature, are strongly changed by radiative corrections associated with the bremsstrahlung emission of soft photons by the initial electrons (see, e. g., [14, 15]). Making allowance for these corrections in a way similar that in Refs (3, 4], for the quantity

$$\Delta R_{t\bar{t}}(s) = \frac{\Delta\sigma_{e^+e^-\to t\bar{t}}(s)}{4\pi\alpha^2/3s}, \tag{20}$$

we get

$$\Delta R_{t\bar{t}}(s) = \frac{18\pi}{m_t^2}\left(Q_t^2 + \frac{v_t^2}{\varkappa(1-m_Z^2/s)^2}\right)\times$$

$$\times \frac{1+3/4\beta_e}{(1-\mathscr{P}(4m_t^2))^2}\int_0^1 dx\, x^{\beta_e-1}\beta_e \operatorname{Im}\Delta G_{E+i\Gamma_t-m_t x}(0,0), \tag{21}$$

where

$$\beta_e = \frac{4\alpha}{\pi}\left(\ln\frac{2m_t}{m_e} - \frac{1}{2}\right),$$

$\mathscr{P}(s)$ is the real part of the polarization operator of the photon.

5. The numerical calculations of $\Delta R_{t\bar{t}}$ were performed according to formula (21) with the following parameters:

$$Q_t = \frac{2}{3}; \quad v_t = 1 - \frac{8}{3}\sin^2\theta_W; \quad \varkappa = (2\sin 2\theta_W)^4;$$

$$\sin^2\theta_W = 0.23; \quad m_Z = 92\ \text{GeV}; \quad \mathscr{P}(4m_t^2) = 0.07.$$

The width Γ_t was calculated by means of the formula (see, e. g., [2]):

$$\Gamma_t = \frac{G_F m_t^3}{8\pi\sqrt{2}}\left(\frac{2P_W}{m_t}\right)\left\{\left(1-\frac{m_b^2}{m_t^2}\right)^2 + \frac{m_W^2}{m_t^2}\left(1+\frac{m_b^2}{m_t^2}\right) - \frac{2m_W^4}{m_t^4}\right\}, \tag{22}$$

where

$$P_W = \frac{[m_t^2-(m_W+m_b)^2]^{1/2}[m_t^2-(m_W-m_b)^2]^{1/2}}{2m_t}. \tag{23}$$

The W-boson and b-quark masses were taken equal to $m_W = m_Z\cos\theta_W = 81$ GeV, and $m_b = 5$, respectively. For $\Delta G_E(0,0)$, the analytical expression (18) with

$$\eta = \pi^2/18\,(0.012)\ \text{GeV}^4 = 0.64\cdot 10^{-2}\ \text{GeV}^4$$

was used. Taking into account the effects of running α_s was simulated as it was done in [3, 4], by using, in (18),

$$\alpha_s = \frac{4\pi}{b_0 \ln\left(\frac{m_t\sqrt{E^2+\Gamma_t^2}}{\Lambda}\right)} \tag{24}$$

with $b_0 = 7.67$, $\Lambda = 100$ MeV.

Figures 1—3 show the values of the ratio $T = \Delta R_{t\bar{t}}/R_{t\bar{t}}$ as a function of energy E for the masses of a t-quark equal to 95, 100 and 110 GeV. In all these cases $R_{t\bar{t}}$ was taken from [3].

These Figures demonstrate a sharp dependence of T on the energy in the vicinity bound states. There is no difficulty in understanding this dependence from a quantitative point of view. For simplicity, let us omit the radiation correction connected with the bremsstrahlung emission of soft photons by the initial electrons. Then

$$T = \frac{\operatorname{Im}\Delta G_{E+i\Gamma_t}(0,0)}{\operatorname{Im} G_{E+i\Gamma_t}(0,0)}, \tag{25}$$

where

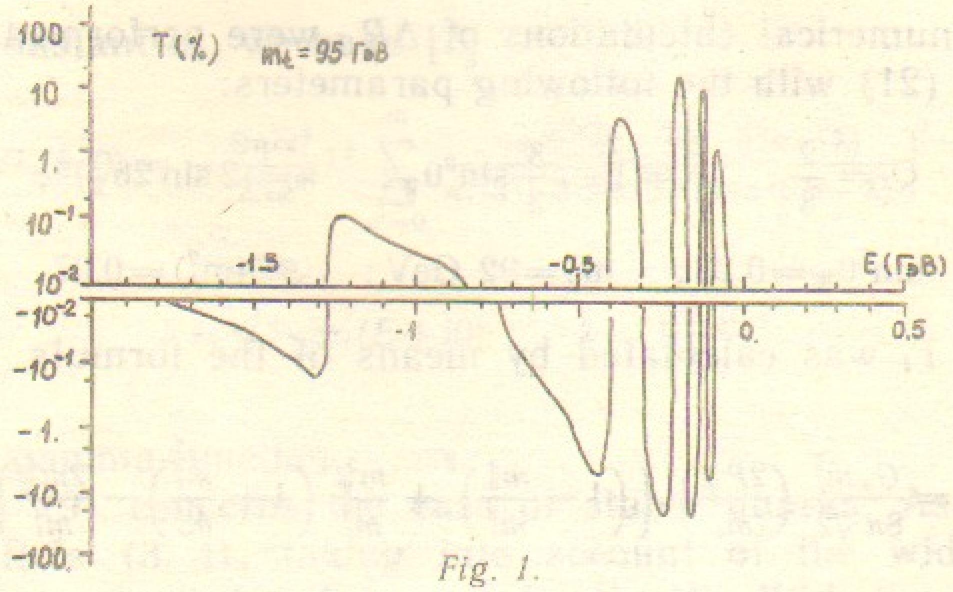

Fig. 1.

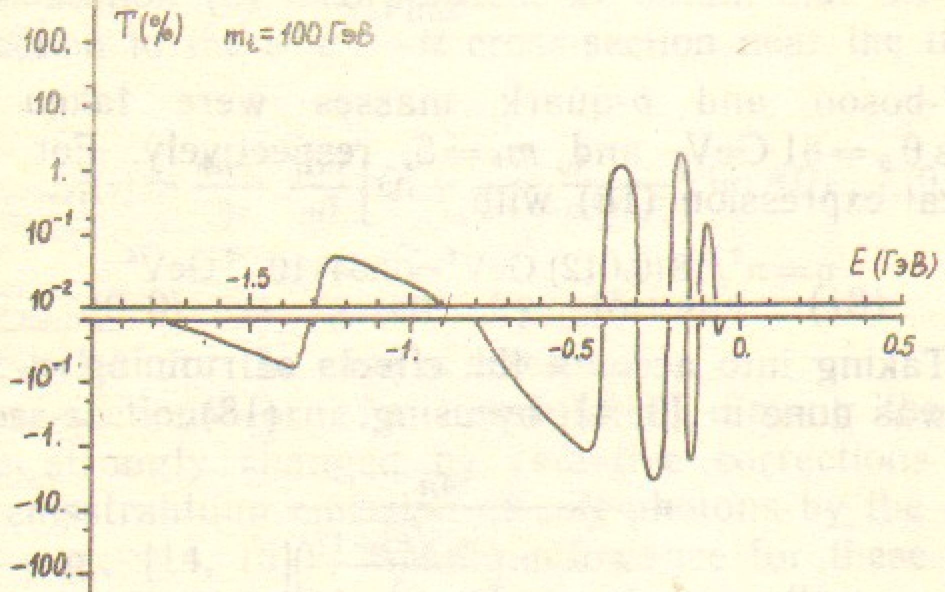

Fig. 2.

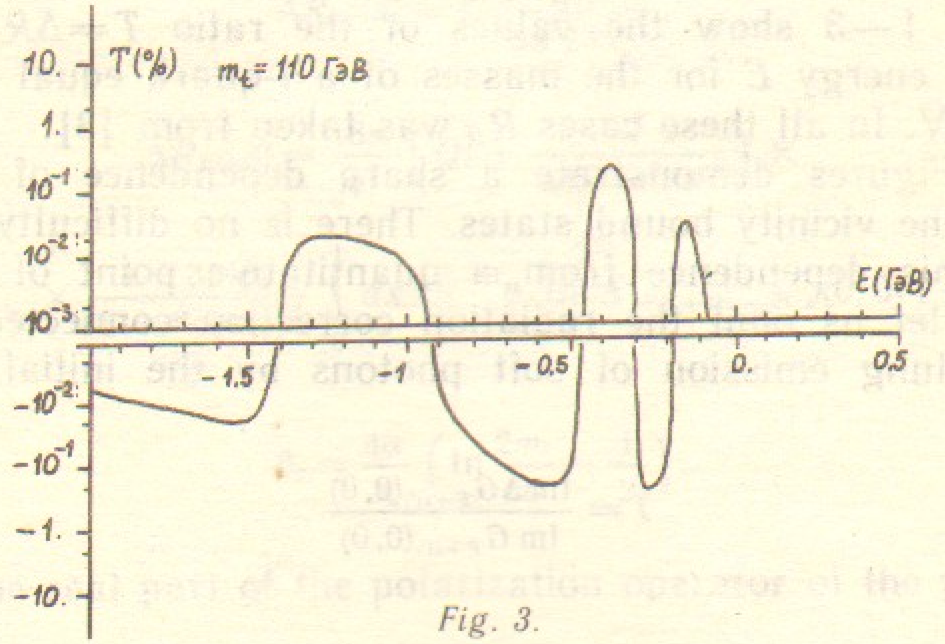

Fig. 3.

$$G_{E+i\Gamma_t}(\vec{r},\vec{r}\,') = \langle \vec{r}\,| \frac{1}{\hat{H}_s - E - i\Gamma_t} |\vec{r}\,'\rangle. \tag{26}$$

We will consider T at E close to the energy of the bound state E_n. In the case when this state is not overlapped with the nearest states due to the width of these states, equal to $2\Gamma_t$, the main contribution to T is determined by this state. Using formulae (16) and (26), we obtain the estimates

$$\Delta G_{E+i\Gamma_t}(0,0) \sim -\eta \frac{|\Psi_n(0)|^2}{(E_n - E - i\Gamma_t)^2} \langle n|\,\vec{r}\,\frac{1}{\hat{H}_8 - E}\,\vec{r}\,|n\rangle \sim$$

$$\sim -\eta \frac{|\Psi_n(0)|^2}{(E_n - E - i\Gamma_t)^2} \frac{r_n^2}{|E_n|}, \tag{27}$$

$$G_{E+i\Gamma_t}(0,0) \sim \frac{|\Psi_n(0)|^2}{E_n - E - i\Gamma_t}. \tag{28}$$

Here $\Psi_n(\vec{r})$ is the wave function of the state in question and r_n is the characteristic distance for n-th level. We have taken into consideration in (27) that Hamiltonian $\hat{H}_8$ corresponds to the Coulomb repulson and, hence, is positively defined. Using (27) and (28), we have

$$T \sim -\frac{\eta r_n^2 (E_n - E)}{|E_n|\,((E_n - E)^2 + \Gamma_t^2)}. \tag{29}$$

Formula (29) gives a qualitatively correct description of the dependence of T on E near E_n at $|E_n| > 2\Gamma_t$:

For $E \simeq E_n = -4m_t\alpha_s^2/9n^2$, the quantity $T(E)$ goes through zero with the positive derivative.

On the right and on the left of E_n at $\Delta E \sim \Gamma_t$ the $|T|$ has a maximum. Since $r_n \sim n^2/m_t\alpha_s$ the magnitude of this maximum grows as the sixth degree of n:

$$|T_n| \sim \frac{\eta n^6}{m_t^3 \Gamma_t \alpha_s^4}. \tag{30}$$

Bearing in mind that Γ_t grows with m_t not slower than according to the asymptotical law (1), we see from formula (30) that $|T|$ falls off with increasing m_t at least as m_t^{-6}. For m_t close to $m_W + m_b$, this decrease occurs much more rapidly.

At $|E_n| < 2\Gamma_t$ the bound states are overlapped, therefore within $|E| < 2\Gamma_t$ the contributions of one and the same order to T give several states, the contributions of different states damping each other as a result. As a consequence, $|T|$ proves to be small.

An additional reduction of $|T|$ is caused by «smearing» due to the emission of soft photons by the initial electrons.

6. As the numerical calculations show, the nonperturbative correction does not exceed 20% at $m_t=95$ GeV, 3% at $m_t=100$ GeV and 0.4% at $m_t=110$ GeV. Note that for each mass the correction reaches a maximum at $-E\sim 2\Gamma_t$. Since at $m_t=95$ GeV the width $\Gamma_t\simeq 0.05$ GeV, the constant chromoelectric field approximation within $-E\sim 2\Gamma_t$ is very doubtful for this mass of the t-quark. Additional errors are in due to the use of the first nonvanishing approximation of a perturbation theory on this field. In view of this, the correction derived should be regarded here only as an estimation. Following this line of analysis, we have no possibilities of improving the result. We may «correct» the perturbation theory near bound states by taking the level shift into account, but the multiple expansion is a fundamental point of the approach and it is impossible to omit it.

The calculations, thus, show that, first, the nonperturbative correction is well pronounced only when its calculation is doubtful and, second, at $m_t > 100$ GeV the correction is insignificant ($<3\%$) and rapidly reduces with m_t. On the one hand, this gives little hope to get information of nonperturbative physics (in particular, on the magnitude of the gluon condensate) from the analysis of the $e^+e^-\to t\bar{t}$ cross-section, but, on the other hand, this offers the possibility of measuring the mass m_t of the t-quark, its width Γ_t and the coupling constant α_s since the cross-section is a calculable function of m_t, Γ_t and α_s [3, 4].

It is not worthless to note that so far we have supposed that the nonperturbative interaction between the gluon field of vacuum fluctuations and b-quarks, produced in the t-decay, may be neglected. Indeed, since b-quarks are produced with relativistic velocities (with respect to each other as well as to t-quarks), this interaction has negligible effect on the cross-section. The usual dimensional estimate gives that the correction to the cross-section is $\sim\langle\frac{\alpha_s}{\pi}(G^a_{\mu\nu})^2\rangle/p_b^4$, where p_b is the typical momenta of the b-quark, $p_b\sim m_t-m_W-m_b$.

When this paper was completed we received the preprint by M.J. Strassler and M.E. Peskin (SLAC-PUB-5308, September, 1990) where the problem of nonperturbative effects is considered

using another approach (in terms of the potential model). The conclusions of the preprint are consistent with ours.

The authors thank V.N. Baier, A.I. Milstein, Yu.F. Pinelis and V.A. Khoze for the interest in the work and numerous discussions.

REFERENCES

1. *Dydak F.* Plenary Talk at the XXIV International Conference on High Energy Physics, Singapore, August 1990.
2. *Bigi I., Dokshitzer Yu., Khoze V., Kühn J., Zerwas P.* Phys. Lett., 1986, v.B181, p.157.
3. *Fadin V.S., Khoze V.A.* Pis'ma v ZhETF, 1987, v.46, p.417; Yad. Fiz., 1988, v.48, p.487.
4. *Fadin V.S., Khoze V.A.* Physics of Elementary Particles. Proc. of XXIV LNPI Winter School, Leningrad, 1989, p.3.
5. *Fadin V.S., Khoze V.A., Sjöstrand T.* Preprint CERN-TH 5687/90, 1990.
6. *Vainstein A.I., Zakharov V.I., Shifman M.A.* Pis'ma v ZhETF, 1978, v.27, p.60.
7. *Voloshin M.B.* Nucl. Phys. B., 1979, v.154, p.365; Yad. Fiz., 1982, v.36, p.247.
8. *Leutwyler H.* Phys. Lett., 1981, v.1398, p.447.
9. *Voloshin M.B., Zaitsev Yu.M.* Usp. Fiz. Nauk, 1987, v.152, p.361.
10. *Baier V.N., Pinelis Yu.F.* Phys. Lett. B., 1982, v.116, p.179.
11. *Poggio E.C., Quinn H.R., Weinberg S.* Phys Rev. D., 1976, v.13, p.1958.
12. *Braun M.A.* ZhETF, 1968, v.54, p.1220.
13. *Barbieri R., Christillin P., Remiddi E.* Phys. Rev. A., 1973, v.8, p.2226.
14. *Baier V.N., Fadin V.S.* Phys. Lett. B., 1968, v.27, p.233.
15. *Lipatov L.N., Khoze V.A.* Proceedings of X LNPI Winter School on Nuclear and Elementary Particle Physics, Part II, Leningrad, 1975, p.409.

V.S. Fadin and O.I. Yakovlev

Nonperturbative Correction to the Threshold Production of $t\bar{t}$-Pairs

В.С. Фадин, О.И. Яковлев

Непертурбативная поправка к сечению образования пары $t\bar{t}$-кварков вблизи порога

Ответственный за выпуск С.Г.Попов

Работа поступила 24 октября 1990 г.
Подписано в печать 30.11 1990 г.
Формат бумаги 60×90 1/16 Объем 1,1 печ.л., 0,9 уч.-изд.л.
Тираж 250 экз. Бесплатно. Заказ № 138

Набрано в автоматизированной системе на базе фотонаборного автомата ФА1000 и ЭВМ «Электроника» и отпечатано на ротапринте Института ядерной физики СО АН СССР,
Новосибирск, 630090, пр. академика Лаврентьева, 11.