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# SEMIHARD HADRON PROCESSES AND QUARK-GLUON STRING MODEL[1]

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# 1. Introduction

As is known the quark-gluon strings model (QGSM) [1-6] based on the $1/N$ expansion in QCD [7-9] had a significant success by the description of different hadron characteristics. For example, it reproduces a large amount of existing experimental data about hadron production at high energies [1-6]. The QGSM has been successfully applied to the production of hadrons containing light $u$, $d$ and $s$ quarks. This model reproduces in great detail and in terms of very few parameters a large amount of existing experimental data about hadron production at high energies in hadron-hadron and hadron-nucleus collisions [1-3]. Besides QGSM describes adequately the production of charmed mesons at the explored low and moderate energies.

However the characteristics integrated over transverse momentum $p_t$ or at average transverse momentum $< p_t >$ are considered only in the framework of this model. So this model is limited usually by the analysis of soft hadron reactions or processes at small transverse momenta $p_t$ of produced hadrons. There are some versions of QGSM taking into account transverse momenta of quarks in the initial hadrons [10, 11, 12] and [13]. It allowed to describe inclusive spectra of hadrons produced in $hN$ collisions up to $p_t = 1. - 1.5$ GeV/c. But there was a large sensitivity of these ones to the initial quark distribution over transverse momentum in a hadron which has been parametrised usually in an exponential form. On the other hand at large $p_t$ the ordinary perturbative QCD can be applied to the analysis of hard processes. Semihard ones can be explained by QCD models of type as [14, 15],etc. However the question arises whether one can apply the QGSM to the analysis of semihard processes if to take into account the colour interaction between colliding quarks (antiquarks) and diquarks (quarks) before the creation and dacay of $q\bar{q}$ (or $q(qq)$) string.

So the phenomenological Pomeron exchange corresponding to the cylinder graph can be understood in the framework of QGSM [1-3] as the exchange of two gluons [16, 17]. However the perturbative calculation of the elastic quark-quark scattering amplitude through two-gluon exchange shows a singularity at $t = 0$. Although this singularity can be cancelled when incorporating quarks in the proton wave function, this procedure is not able to reproduce the $t$-dependence of the differential cross section observed in experiment [18]. Meanwhile some time ago the Pomeron was described by the exchange of two nonpertubative gluons [19, 20, 21], where by nonpertubative it means a gluon whose propagator does not show a pole at $k^2 = 0$. In [22] the two-gluon exchange model was applied to the description of nucleon-nucleon scattering where a gluon propagator regularized by a dynamically generated gluon mass, derived by Cornwall some years ago [23].

The production of hadrons at high energies is described in the QGSM by "cutting" the forward scattering diagrams of the "cylindrical" type [1-3]. Each cylinder corresponds to the exchange of a single Pomeron. Such Pomeron exchange as the one of two non-perturbative gluons can be considered for the calculation of cylinder-cut graphs in the framework of the QGSM. This paper is dedicated to the exploitation of this idea to the analysis of hadron processes in the framework of QGSM.

## 2. General formalism

Consider the inclusive spectrum of hadrons produced in nucleon-nucleon collision in the framework of QGSM. As is known the main contribution to such processes at large energies gives the graphs of cylinder-cut type in the $s$ channel corresponding to the Pomeron exchange in $t$ channel [1-6]. The hadron production can be considered in the following manner [11]. Each of two colliding nucleons is divided into a quark and a diquark with the opposite transverse momenta. After the colour interaction (by means of non-perturbative gluon) between the quark of first proton and diquark of another one, and diquark of the first proton and quark of second one two quark-gluon strings are created in the chromostatic constant field; then these two strings decay into secondary hadrons. This process is repeated $n$ times during the



production of $n$ Pomeron showers (or 2n $q\bar{q}$ chains). Such division of the transverse momentum considered in [11] is analogous to the so-called consequent division of energy between $n$ Pomeron showers in the multiperipheral model of the hadron production [2].

In principle, there can be so-called regular division of energy and transverse momentum between 2n $q\bar{q}$ chains or $n$ Pomeron showers. But at all versions of QGSM nobody considered the colour interaction between quark an antiquark or diquark and quark respectively. For the soft processes, i.e., at small $p_t$, quark distributions over $x$ on the ends of $q - \bar{q}$ strings are found from Regge asymptotic of graphs (cylinder or planar type) at $x \to 1$ [1-5]. The decay of each this string into hadrons is described by fragmentation functions, the form of which is determined by the Regge asymptotic of these graphs too [1-5]. The dependence of hadron inclusive spectra on $p_t$ is the result of the inclusion of internal hadron transverse momenta of quarks only [10, 11] and [13]. This approach is good to be applied to hadron production at $p_t \leq 1. - 1.5$ GeV/c [10, 11, 12].

However if we want to analyze the hadron production processes at not large $p_t$, for example, $p_t \leq 3 - 4$ Gev/c, the question arises whether one can use the QGSM for this aim. In order to do it we must know the distribution of quarks (antiquarks) and diquarks over transverse momenta $k_t$ on the ends of $q\bar{q}$ string after some colour interaction between constituents of colliding hadrons. For this aim one can use the approach suggested in [22] where the Pomeron is described by the exchange of two nonperturbative gluons and it means that a gluon propagator has not a pole at $k^2 = 0$ [20]. For example, in [21] an approximate Shwinger-Dayson [SD] equation for the gluon propagator in the axial gauge has been solved and found a solution whose infrared behaviour is less singular than a pole at $k^2 = 0$. This approach has given reasonable agreement with experimental data. Recently in [22] it has been shown how a gluon propagator regularized by a dynamically generated gluon mass, derived in [23] some years ago, successfully describes nucleon-nucleon scattering in two-gluon exchange model. The gluon propagator represented as an approximate solution of the SD equation and it had not the infrared singularity at $k^2 = 0$ as a result of some cut-off parameter $m_0$ or so-called "effective" gluon mass. In fact the model of ref. [22] has a single parameter which was taken to be the ratio of $m_0$ and $\Lambda = \Lambda_{QCD}$.

In this paper we take into account the interaction between quark and antiquark (or diquark and quark) before the creation and the decay of the corresponding string. The main contribution to this one results in the one gluon-exchange between them, which can be calculated using the gluon propagator with the cut-off parameter as like as in [22, 23]. So this interaction can be considered as the elastic $q\bar{q}$ or quark-diquark ($q(qq)$) scattering calculated in the one-gluon exchange approximation when transverse momenta of quarks or diquarks are changed only. If one represents the distribution of quarks or diquarks in a nucleon over $x$ and $p_t$ in the factorized form, i.e.

$$f_\tau(x,\vec{k}_t) = f_\tau(x) g_\tau^{(0)}(\vec{k}_t), \tag{1}$$

where $\tau$ means the flavour of a quark, $x$ and $\vec{k}_t$ are its fraction of the longitudinal momentum and the transverse one respectively; $f_\tau(x)$ is the $x$-distribution and $g_\tau^{(0)}(\vec{k}_t)$ is the $k_t$-distribution of this quark (antiquark or diquark) in the initial hadron. Then the factorized form of $f_\tau(x,\vec{k}_t)$ will be true after the one-gluon exchange, but the distribution $g_\tau^{(0)}(\vec{k}_t)$ will be changed only. After such interaction this one can be calculated by the following manner:

$$g_\tau^{(1)}(\vec{k}_{1,t}) = C_1 \int D^2(q^2) g_\tau^{(0)}(\vec{k}_t)\delta^{(2)}(\vec{k}_{1,t} - \vec{k}_t - \vec{q}_t) d^2q_t\, d^2k_t \tag{2}$$

or

$$g_\tau^{(1)}(\vec{k}_{1,t}) = C_1 \int D^2[(\vec{k}_{1,t} - \vec{k}_t)^2] g_\tau^{(0)}(\vec{k}_t) d^2k_t\,.$$

where $D(q^2)$ is the regularized gluon propagator and according to [22, 23] it can be written in the following form:

$$D(q^2) = \frac{\alpha_s(q^2)}{q^2 + m^2(q^2)}\,. \tag{3}$$

where:

$$\alpha_s(q^2) = \frac{1}{b_0 ln \frac{q^2+4m^2(q^2)}{\Lambda^2}}, \tag{4}$$

$$m^2(q^2) = m_0^2 \left[\frac{ln((q^2+4m_0^2)/\Lambda^2)}{ln(4m_0^2/\Lambda^2)}\right]^{-12/11}. \tag{5}$$

Here in (4), (5) $b_0 = (33 - 2n_f)/48\pi^2$; $m_0 = 0.37$ Gev for $\Lambda = 0.3$ GeV. As in our previous paper [11, 12] we choose the quark $p_t$-distribution in the initial hadron in the Gauss form normalized to 1,

$$g_\tau^{\cdot}(\vec{k}_t) = \frac{\gamma}{\pi} exp(-\gamma \vec{k}_t^2). \tag{6}$$

with $\gamma$ to be a free parameter, $\gamma = 1/ < k_t^2 >$ where $< k_t^2 >$ is the average value of the squered transverse momentum of quarks in a nucleon. The constant $C_1$ in (2) is determined by the normalization condition:

$$\int g_\tau^{(1)}(\vec{k}_t) d^2k_t = 1. \tag{7}$$



As is seen from (3) the gluon propagator, $D(q^2)$, at very large $q^2$ returns to the usual one asymptotically, i.e. $D(q^2) \propto q^{-2}$, because $m^2(q^2) \to 0$ at $q^2 \to \infty$. This method of the calculation of the $k_t$-distribution of a quark (antiquark or a diquark) on the ends of $q\bar{q}$ string allows us to construct inclusive spectra of hadrons produced from its decay as the function of $x$ and $p_t$. Note the sensitivity of $g_\tau^{(1)}$ calculated using (2) to the choice of the $g_\tau^{(0)}$ form is very small if we consider not small $k_t$. It is caused by the strong dependence of the gluon propagator (3) on $q_t$.

Let us return to this construction in the framework of the version of QGSM taking into account transverse momenta of quarks [11, 12]. The expression for the invariant inclusive hadron spectrum for the reaction $pp \to hX$ corresponding to graphs of the cylinder-cut type [1-6] can be written in the following form [11]:

$$\rho_{pp\to hX}(x,\vec{p}_t) \equiv E\frac{d\sigma_A}{d^3\vec{p}} = \sum_{n=1}^{\infty} \sigma_n(s)\phi_n(x,\vec{p}_t), \tag{8}$$

where $\sigma_n$ is the cross section for production of the $n$ pomeron chain (or $2n$ quark-antiquark strings) decaying into hadrons, calculated in the frame of the "eiconal approximation" [24], $\phi_n(x,\vec{p}_t)$ is the $x$ and $p_t$-distributions of hadrons produced in the decay of the $n$ pomeron chain. These functions, $\phi_n(x,p_t)$, were represented in the following form [11]:

$$\phi_n(x,\vec{p}_t) = \int_{x_+(n)}^{1} dx_1 \int_{x_-(n)}^{1} dx_2 \Psi_n(x_n,\vec{p}_t;x_1,x_2); \tag{9}$$

with

$$\begin{aligned}\Psi_n(x_n,p_t;x_1,x_2) = {} & F_{qq}^{(n)}(x_+(n),\vec{p}_t;x_1)F_{q_v}^{(n)}(x_-(n),\vec{p}_t;x_2)/\bar{F}_{q_v}^{(n)}(0,\vec{p}_t) + \\ & F_{q_v}^{(n)}(x_+(n),\vec{p}_t;x_1)F_{qq}^{(n)}(x_-(n),\vec{p}_t;x_2)/\bar{F}_{qq}^{(n)}(0,\vec{p}_t) + \\ & 2(n-1)F_{q_{sea}}^{(n)}(x_+(n),\vec{p}_t;x_1)F_{q_{sea}}^{(n)}(x_-(n),\vec{p}_t;x_2)/F_{q_{sea}}^{(n)}(0,p_t)\end{aligned}$$

where $x_\pm = x_\pm(n) = 0.5(\sqrt{x_t^2 + x_n^2} \pm x)$, $x_t = 2\sqrt{(m_h^2 + \vec{p}_t^2)/s}$ ; $m_h$ is the mass of the produced hadron, $s$ is the total energy squared of colliding protons in c.m.s. Note that Feynman variable $x$ in $n$-Pomeron shower depends on $n$ [11]: $x_n = x_/(1-x_0)^{n-1}$, $x_0 \simeq 0.25$.

The functions $F_\tau^{(n)}(x_\pm(n),\vec{p}_t)$ are represented by convolutions

$$F_\tau^{(n)}(x_\pm(n),\vec{p}_t;x_{1,2}) = \int d^2k_t f_\tau^{(n)}(x_{1,2},\vec{k}_t) G_{\tau\to h}(x_\pm(n)/x_{1,2},\vec{k}_t;\vec{p}_t), \tag{10}$$

$$F_\tau^{(n)}(0,\vec{p}_t;x_{1,2}) = \int d^2k_t f_\tau^{(n)}(x_{1,2},\vec{k}_t) G_{\tau\to h}(0,\vec{p}_t) = G_{\tau\to h}(0,\vec{p}_t). \tag{11}$$

Here $\tau$ means the flavour of the valence (or sea quark) or diquark, $f_\tau^{(n)}(x,\vec{k}_t)$ is the quark distribution function, depending on the longitudinal momentum fraction $x$ and the transverse momentum $\vec{k}_t$ in the $n$-Pomeron chain:

$$G_{\tau\to h}(z,\vec{k}_t;\vec{p}_t) = zD_{\tau\to h}(z,\vec{k}_t;\vec{p}_t), \tag{12}$$

$D_{\tau\to h}(z,\vec{k}_t;\vec{p}_t)$ is the fragmentation function of a quark or diquark of flavour $\tau$ into a hâdron $h$.

The functions $f_\tau^{(n)}(x)$ and $g_\tau^{(n)}(\vec{k}_t)$ can be calculated by using regular [10] or consequential [11] division of $x$ and $k_t$ between 2n $q\bar{q}$ chains. Here we use the second one [11], so the $g_\tau^{(n)}(\vec{k}_t)$ are calculated according to [11] by the following manner. The quark functions $f_\tau(x,\vec{k}_t)$ in initial hadrons are represented in the factorized form (1). The distribution functions over transverse momentum $p_t$, $g_\tau^{(n)}(\vec{k}_{n,t})$, are calculated analogous to (2). This means that the quark distribution after one nonperturbative gluon exchange is given by (2). After exchange by the second gluon the quark distribution in corresponding chain will be expressed already via the function $g_\tau^{(1)}(\vec{k}_t)$, i.e.:

$$g_\tau^{(2)}(\vec{k}_{2,t}) = C_2\int g_\tau^{(1)}(\vec{k}_{1,t})D^2[(\vec{k}_{2,t}-\vec{k}_{1,t})^2]d^2k_{1,t}. \tag{13}$$

Repeating this iteration procedure we obtain the quark distribution function in the $n$ chain expressed via the function $g_\tau^{(n-1)}(\vec{k}_{n-1,t})$ and therefore via the function $g_\tau^0(\vec{k}_t)$:

$$\begin{aligned} g_\tau^{(n)}(\vec{k}_{n,t}) &= C_n\int g_\tau^{(n-1)}(\vec{k}_{n-1,t})D^2[(\vec{k}_{n,t}-\vec{k}_{n-1,t})^2]d^2k_{n-1,t} = \\ &\int d^2k_{n-1,t}D^2[(\vec{k}_{n,t}-\vec{k}_{n-1,t})^2]\int d^2k_{n-2,t}D^2[(\vec{k}_{n-1,t}-\vec{k}_{n-2,t})^2]\ldots \\ &\cdot\int d^2k_tD^2[(\vec{k}_{1,t}-\vec{k}_t)^2]g_\tau^{(0)}(\vec{k}_t). \end{aligned} \tag{14}$$

It is easily to see that in the $n$-pomeron chain the quark functions will be factorized too:

$$\tilde{f}_\tau^{(n)}(x_n,\vec{k}_{n,t}) = f_\tau^{(n)}(x_n)g_\tau^{(n)}(\vec{k}_{n,t}). \tag{15}$$

This calculation of $k_t$-distribution of quarks on the ends of $2n$-th sring is a principal difference from our previous papers [11, 12]. But $x$-distribution after $n$-Pomeron exchanges $f_\tau^n(x)$ is calculated according to [11], i.e. the Feynman variable $x$ in $n$-Pomeron shower depends on $n$ (see above).

The functions $\tilde{G}_{\tau\to h}(z,\vec{k}_t;\vec{p}_t)$ have represented in the following form [11]:

$$\tilde{G}_{\tau\to h}(z,\vec{k}_t;\vec{p}_t) = G_{\tau\to h}(z,\vec{p}_t)\tilde{g}_{\tau\to h}(\tilde{k}_t), \tag{16}$$



where

$$\tilde{g}_{\tau\to h}(\tilde{k}_t) = \frac{\tilde{\gamma}}{\pi}exp(-\tilde{\gamma}\tilde{k}_t^2), \tag{17}$$

$$\tilde{k}_t = \vec{p}_t - z\vec{k}_t, \qquad z = \frac{x_\pm}{x_{1,2}}. \tag{18}$$

Insert now the functions (14)-(17) to (10) and integrate it over $d^2k_t$. In principle, it can be made analytically if the function $g_\tau^{(n)}(\vec{k}_t)$ calculated by (10) is approximated by a sum of exponentials over $k_t^2$ (see Appendix).

## 3. Results and discussions

The results of calculation of differential cross sections performed with the help of the modified QGSM are presented in Figs. 1-5. Using (8) we can calculate different characteristics of hadrons produced in N-N collisions, for example, correlation of the average transverse momentum $< p_t >$ and hadron multiplicity $N$ (see for example [11]). Besides the inclusive spectra of hadrons defined by (8), we can calculate the distributions of hadrons over $x$ or $p_t$,i.e.:

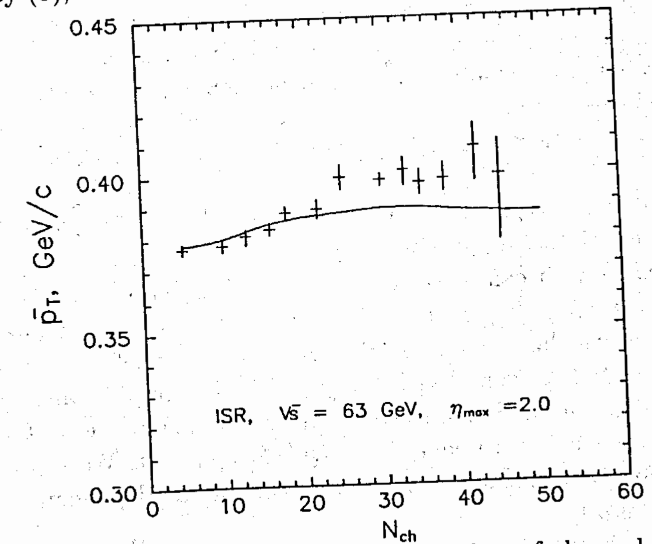


Fig. 1. Dependence of $< p_t >$ on the number of charged particles, $N_{ch}$. The curve shows our calculation for $\sqrt{s} = 63$ GeV. Symbols denote the experimental data for $\sqrt{s} = 63$ GeV, rapidity interval, $|y| = 2$ [25].

$$F(x) = \int \rho_{pp\to hX}(x, p_t) d^2 p_t \tag{19}$$

and

$$\frac{d\sigma}{dp_t^2} = \frac{\pi}{2}\sqrt{s}\int \varrho_{pp\to hX}(x,p_t)\frac{dx}{E^*} \tag{20}$$

where $E^*$ is the energy of final hadrons in N-N c.m.s.

The results of the calculations of $< p_t > (N)$ for $pp$ collisions are shown in Fig. 1. The calculations reproduce the data at small multiplicities $N$ and go below at large $N$. The discrepancy between the results of the calculations and the experimental data about $< p_t > (N)$ may be connected with the fact that the expressions for the cross sections $\sigma_n$ (see (8)) were obtained with inclusion of only non-enhanced graphs of Reggeon theory [2]. Inclusion of enhanced-type diagrams leads to the appearance of $1/x$ terms in the distributions for sea quarks, whose contribution is especially large for $x \simeq 0$.

The calculation results for inclusive hadron spectra performed using the formulae (8), (20) are presented in Figs. 2-5. In Figs. 2, 3 we show the inclusive spectra of $\pi^+$ and $\pi^-$ mesons produced in $pp$ collisions at $p_{lab} = 200$ GeV/c as a function of $p_t$ for three values of the Feynman variable x: $x = 0$, 0.3 and 0.6. We see a good description of experimental data with modified QGSM. Corresponding parameter values in formulas (6) and (17) are: $\gamma = 9\,(\mathrm{GeV/c})^{-2}$ what corresponds to $< k_t^2 >= 1/\gamma \simeq 0.1\,(\mathrm{GeV/c})^2$, $\tilde{\gamma} = 7\,\mathrm{GeV}^{-2}$. Note that the calculation results depend on the value of this parameters very weakly.

In Fig. 4 we show the invariant cross section $E(d\sigma/d^3\vec{p}$ versus transverse momentum $p_t$ for the $\pi^0$ for mesons produced in $pp$ collisions at $\sqrt{s} = 52.7$ GeV. From this figure it can be seen that our version of the QGSM gives good description of the data [27] up to $p_t < 3$ GeV/c. Some discrepancy with data at larger $p_t$ can be caused by the following. First of all the $p_t$ quark distribution in initial hadrons has been taken in the simple Gauss form (6) which is true at small $p_t$. Second reason can be conjugated with the Regge trajectories, $\alpha(t)$, which are taken in the QGSM at $t = 0$. We suppose that in the calculations it is needed to take into account the dependence $\alpha(t)$ on transfer momentum $t$.

In Fig. 5 we present the calculation results for the differential cross sections $d\sigma/dp_t^2$ of the reactions $pp \to D^n X$, $n = +(a), -(b), 0\,(c), \bar{0}\,(d)$, at $\sqrt{s} = 27.4$ GeV. The curves correspond to the case when the Regge $\Psi$- trajectory has intercept, $\alpha_\Psi(O) = 0$, which corresponds to the nonlinear $\Psi$-trajectory.



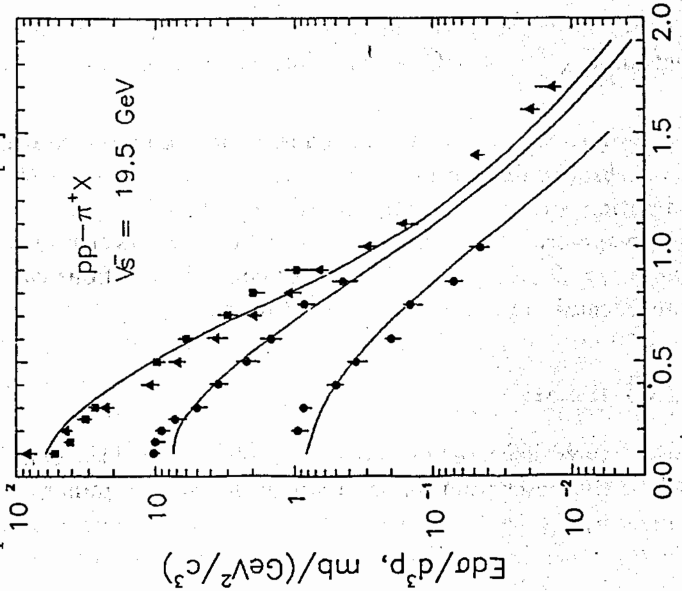


Fig. 2. Invariant cross section of $\pi^+$ mesons produced in reaction $pp \to \pi^+ X$ at $p_{lab} = 200$ GeV/c as a function of $p_t$ for the values $x = 0$, 0.3 and 0.6. Experimental data are from Refs. [26].

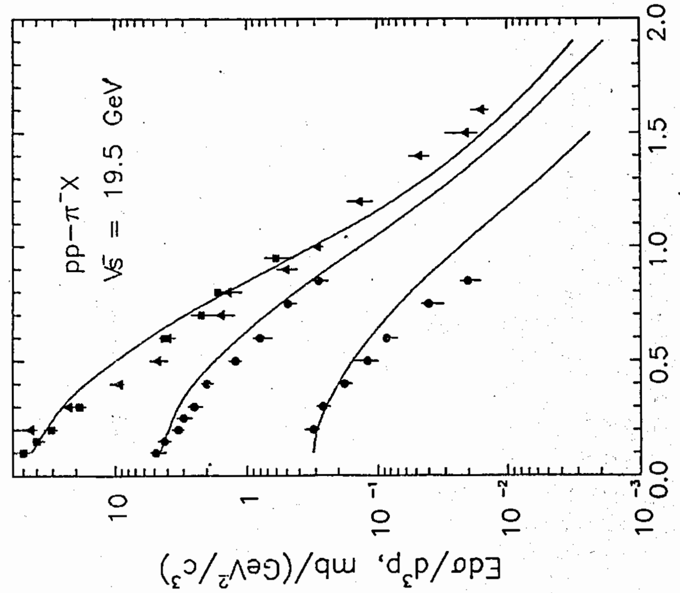


Fig. 3. Invariant cross section of $\pi^-$ mesons produced in reaction $pp \to \pi^- X$ at $p_{lab} = 200$ GeV/c as a function of $p_t$ for the values $x = 0$, 0.3 and 0.6. Experimental data are from Refs. [26].

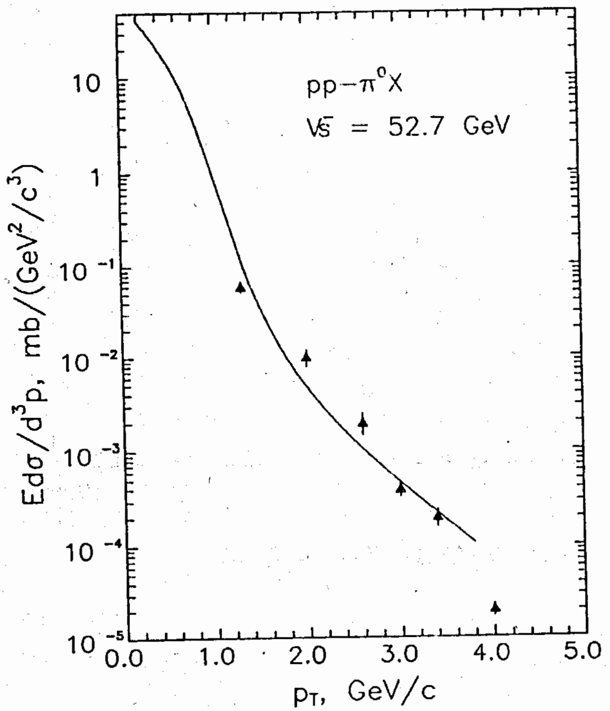


Fig. 4. The inclusive cross section of the reaction $pp \to \pi^0 X$ versus transverse momentum at $\sqrt{s} = 52.7$ GeV. The data are from Ref.[27].

The respective values of parameters in fragmentation functions are [12]: $a_0 = 0.1\,10^{-3}$, $a_1 = 5$. From this figure it can be seen that the version of the QGSM under consideration gives good description of the data [28].

The inclusive spectra of all $D$-mesons have been calculated with the parameter value $\tilde{\gamma}$, $\tilde{\gamma} = 2$ (Gev/c)$^{-2}$. The better coincidence of our calculations with the experimental data is for $D^+$ and $D^0$ mesons.

## 4. Conclusion

In this work we have discussed one possible modification of the Quark-Gluon String Model to the description of semihard processes. Semihard processes can be explained by the QCD models such as [14, 15].etc. The main question considered in this work is whether one can apply the QGSM to the analysis



of semihard processes if one includes the colour interaction between colliding quarks (antiquarks) and diquarks (quarks) before the creation and dacay of $q\bar{q}$ (or $q(qq)$) string.

Our approach to the analysis of soft and semihard hadron processes in the QGSM takes into account the dependence of quark distributions in hadrons and the quark fragmentation functions on the transverse momentum $p_t$. In the framework of the QGSM the colour interaction of valence quarks and diquarks and sea quarks (antiquarks) of colliding hadrons has been taken into account. This one has been calculated as the exchange by one nonpertubative gluon, i.e., the cut-off parameter in the gluonic propagator has been included. The method proposed to include the $p_t$-dependence in the QGSM is close to the successive division of the transverse momentum $p_t$ and $x$ between $n$-Pomeron showers or $2n$ quark-antiquark chains considered in our early works [11, 12]. This method shows a strong dependence of hadron characteristics on $n$.

The QGSM modified by such a way has allowed us to analyze the correlation $< p_t > (N)$ and inclusive hadron spectra produced in hadron collisions at transverse momenta up to $3-4$ GeV/c (Figs. 1-5). The calculations reproduce the data on $< p_t > (N)$ at small multiplicities $N$. To describe the data at large $N$ one needs, as was mentioned above, to take into account the graphs of enhanced type of Reggeon theory [2]. Inclusion of such diagrams results in the appearance of $1/x$ terms in the distributions for sea quarks, which give especially large contribution for $x \simeq 0$. Modified QGSM can reproduce a lot of data about invariant cross sections versus $p_t$, for example of $\pi^\pm$ mesons produced in $pp$ collisions at different values of x.

In the framework of new version of QGSM, the description of the differential cross sections of $D^+$, $D^-$, $D^0$ and $\bar{D}^0$ mesons versus transverse momentum $p_t$ has been performed. The comparison of the calculations with data show good agreement for the case when Regge $\Psi$-trajectory has intercept, $\alpha_\Psi(O) = 0$. Such intercept value corresponds to the nonlinear $\Psi$ trajectory. Main uncertainties of the calculations come from poor knowledge of intercept of $\Psi$-trajectory. At high energies, the predictions for $\alpha_\Psi(0)$ differ by a factor close to 3. Unfortunately , the large error in charm cross section measurement at $\sqrt{s} = 630$ GeV does not allow one to extract a useful constant on $\alpha_\Psi(0)$.

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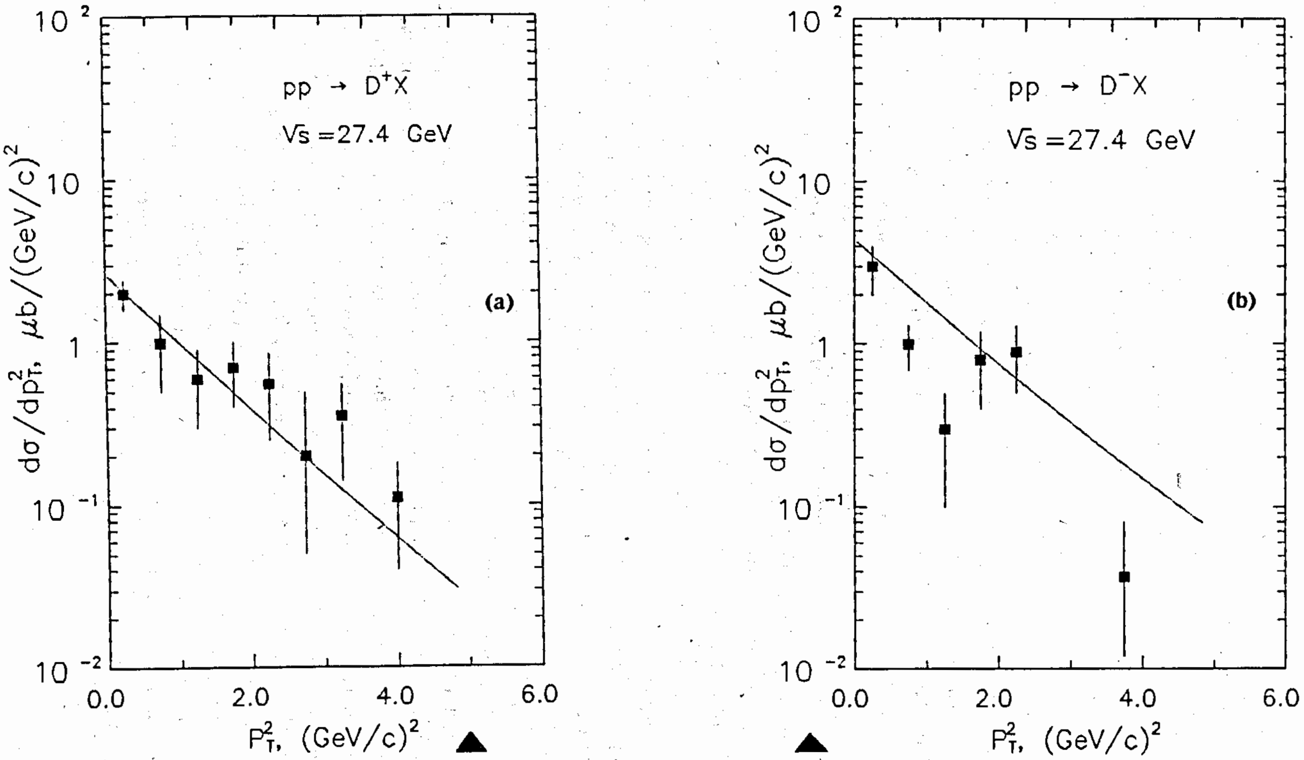
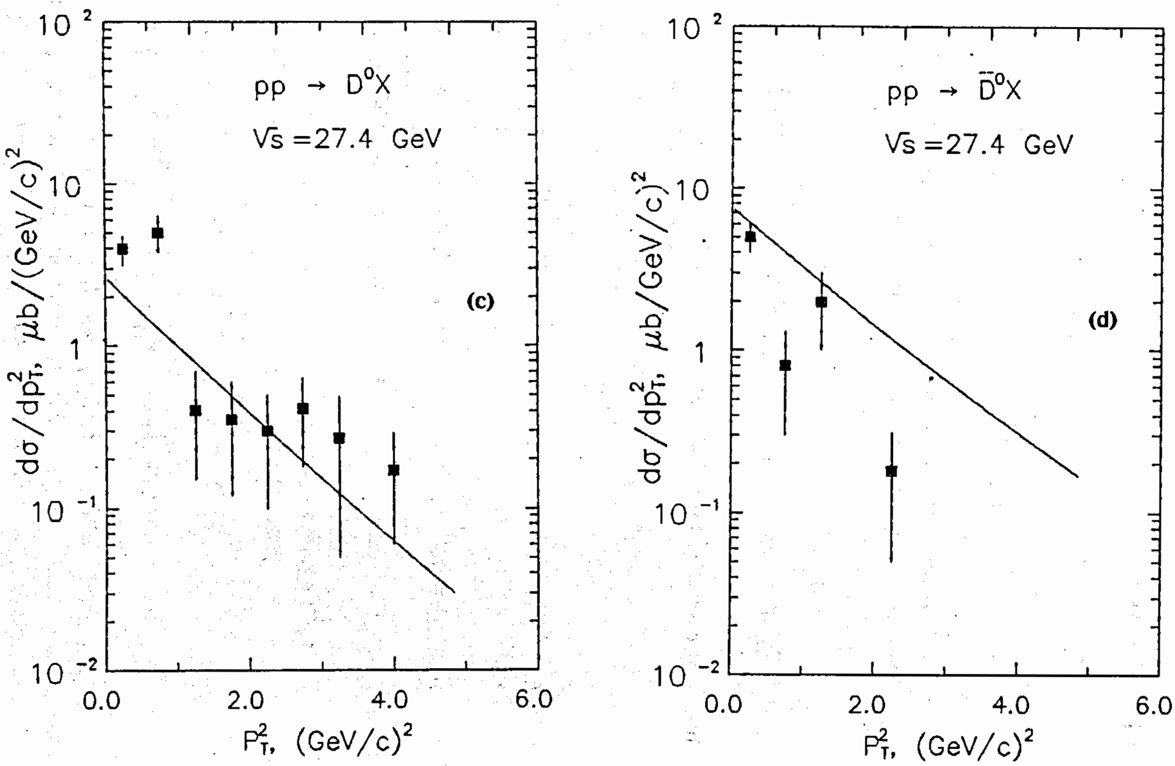



Fig. 5. Differential cross sections of the reactions $pp \to D^+X$ (a), $pp \to D^-X$ (b), $pp \to D^0X$ (c), $pp \to \bar{D}^0X$ (d) at the energy $\sqrt{s} = 27.4$ GeV as a function of the square $p_t^2$ of the transverse momentum. Solid curves correspond to calculations with the intercept $\alpha_\psi(0) = 0$. The data are from Ref. [28].





## 5. Appendix

Show, if the function $D^2(q_t^2)$ can be approximated in form:

$$D^2(q_t^2) = \sum_{i=1}^{N} a_i exp(-b_i q_t^2), \qquad (21)$$

where $a_i, b_i$-some parameters, then the integration of (10) is performed analytically. Consider firstly the simple case $n = 1$, then substituting (21) and (6) to (2) we have:

$$g_\tau^{(1)}(\vec{k}_{1t}) = \frac{\gamma}{\pi}\sum_{i=1}^{N} a_i \int exp(-\gamma k_{1t}) exp(-b_i q_t^2)\delta^{(2)}(\vec{k}_t - \vec{k}_{1t} - \vec{q}_t) d^2k_{1t} d^2q_t. \qquad (22)$$

Integrating (22) over $d^2q_t$ and $d^2k_{1t}$ we get:

$$g_\tau^{(1)}(k_t^2) = \sum_{i=1}^{N} \frac{a_i}{b_i} d_i exp(-d_i k_t^2), \qquad (23)$$

where $d_i = \frac{b_i\gamma}{b_i+\gamma}$ Substituting now (1) and (16) to (10), we get the following expression for $F_\tau^{(1)}(x_\pm, \vec{p}_t; x_{1,2})$ :

$$F_\tau^{(1)}(x_\pm, \vec{p}_t; x_{1,2}) = f_\tau^{(1)}(x_{1,2}) G_{\tau\to h}(z,\vec{p}_t) \int g_\tau^{(1)}(k_t)\tilde{g}_{\tau\to h}(\vec{p}_t - z\vec{k}_t) d^2k_t \cdot \qquad (24)$$

Substituting now (23) and (17) to (24) we integrate it over $d^2k_t$ and finally have:

$$F_\tau^{(1)}(x_\pm, \vec{p}_t; x_{1,2}) = f_\tau^{(1)}(x_{1,2}) G_{\tau\to h}(z,\vec{p}_t) I_1(z,\vec{p}_t), \qquad (25)$$

where:

$$I_1(z,\vec{p}_t) = \sum_{i=1}^{N} \frac{a_i}{b_i}\gamma_{zi} exp(-\gamma_{zi} p_t^2), \qquad (26)$$

here $\gamma_{zi} = \frac{d_i\tilde{\gamma}}{d_i + z^2\tilde{\gamma}}$ But for any n (10) is integrated over $d^2k_t$ anologous to this simple case.

The question arises whether $D^2(q_t^2)$ (see (3)) can be represented in the form (18). This approximation can be true at the following parameters: $N = 5, a_1 = 0.31, a_2 = 0.333, a_3 = 0.172, a_4 = 0.155, a_5 = 0.026$; and $b_1 = 9.294, b_2 = 2.78, b_3 = 3.619, b_4 = 0.882, b_5 = 184$.

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Лыкасов Г.И., Сергеенко М.Н. E2-95-62
Полужесткие адронные процессы и модель кварк-глюонных струн

Предлагается новый подход к анализу мягких и полужестких процессов. В рамках модели кварк-глюонных струн учитывается взаимодействие валентных кварков и дикварков, а также морских кварков (и антикварков). Такое взаимодействие вычисляется как обмен непертурбативным глюоном, т.е. вводится параметр обрезания в глюонном пропагаторе. Такая процедура позволяет нам анализировать инклюзивные спектры адронов в адронных столкновениях при поперечных импульсах вплоть до 3—4 ГэВ/с.

Работа выполнена в Лаборатории ядерных проблем ОИЯИ.

Препринт Объединенного института ядерных исследований. Дубна, 1995

Lykasov G.I., Sergeenko M.N. E2-95-62
Semihard Hadron Processes and Quark-Gluon String Model

A new approach to the analysis of soft semihard hadron processes is suggested. In the frame of the Quark-Gluon String Model the interaction of valence quarks and diquarks and sea quarks 9antiquarks) of colliding hadrons is taken into account. This one is calculated as the exchange by one nonperturbative gluon, i.e., the cut-off parameter in the gluonic propagator is included. This one allows us to analyze the inclusive hadron spectra in hadron collisions at transverse momenta up to 3—4 GeV/c.

The investigation has been performed at the Laboratory of Nuclear Problems, JINR.

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