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DIQUARK DESTRUCTION AND PREASYMPTOTICAL MECHANISMS OF ANTIPROTON ANNIHILATION

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1. Introduction

Impossibility of describing soft hadronic processes, starting from the first QCD principles, leads to appearance of QCD motivated model-dependent approaches. Quite fruitful one is the so-called quark-gluon string model (QGSM) /1-8/, which is based on 1/N ($N=N_c$, N_c/N_f=const) expansion method in QCD /9/. In high energy reactions without baryons the 1/N expansion in the QCD is turned out to be equivalent /9/ to the topological expansion /10, 11/. Reggeon exchange is assumed to be described by analogy with the topological expansion by the planar graphs, but the Pomeron corresponds to the graph of cylindrical topology.

The following space-time picture of hadron production, motivated by the confinement is also assumed. As the result of colour exchange between the colliding particles, the chromoelectric tubes (strings) are emerged /12,13/ and then they fall into pieces. The cuts of planar or cylindrical graphs correspond to hadronization of one or two colour triplet strings respectively. The hadronization process is supposed in the QGSM to be independent on the history of the string origin, but it is only determined by the type of quarks on the string ends.

In practice the QGSM deals with reactions including baryons (for the certainty we restrict ourself by nucleons below), which are considered as a system which consists of quark and structureless object - diquark (D). Since diquark is a colour antitriplet, nucleon can be considered as meson in those reactions where diquark is not destroyed. Such approach describes well /1-8/ a bulk of experimental data.

However $\bar{N}N$ annihilation at high energies is accompanied by the diquark destruction (diquark exchange contribution is small /14/), thus, it is worthy of special discussion. The important fact, which makes it difficult to use the methods of 1/N expansion /15,16/ is: the number of quarks in nucleon is equal to N_c without fail.

Rossi and Veneziano /15/ have proposed to classify annihilation mechanisms not on the ground of 1/N expansion, but using some topological classification. The important innovation of their consideration is the conception of string junction (SJ). The space-time picture of nucleon movement is presented by three sheets, corresponding to three quarks. The line of interception of these planes is trajectory of SJ in space.

The graphs in fig. 1a, b, c correspond to Reggeons containing the SJ-$\overline{SJ}$ pairs except quarks.

Following [15,16], we denote the Reggeons in fig. 1a, b, c by M_4^J, M_2^J, M_0^J respectively. Their cuts correspond to the string configurations, shown in fig. 2, which are the annihilation channels (we neglect here sea $\bar{N}N$ pair production).

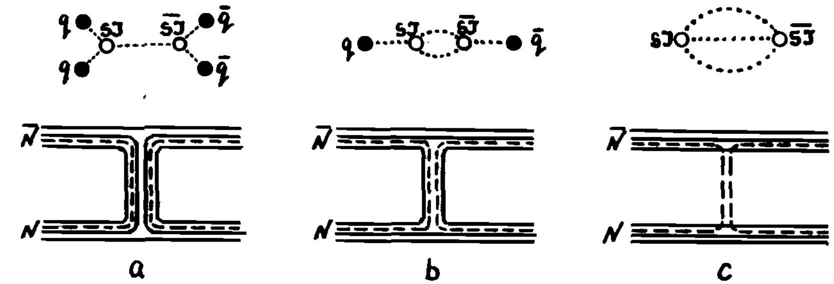

Fig. 1. Graphs with exchange of Reggeons, containing the string junctions. The solid and dashed lines correspond to quarks and to SJ. There are shown on the top of the picture string configurations (dotted lines) of corresponding mesons M_4^J, M_2^J and M_0^J.

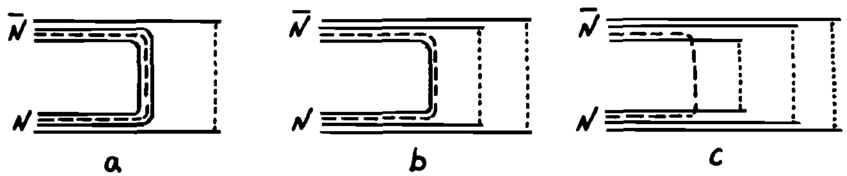

Fig. 2. Processes corresponding to the cuts of the graphs in fig. 1.

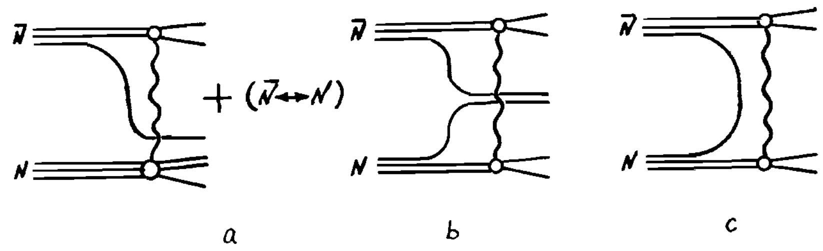

Fig. 3. Perturbative QCD graphs, which are responsible for the annihilation. The waved line corresponds to a gluon.

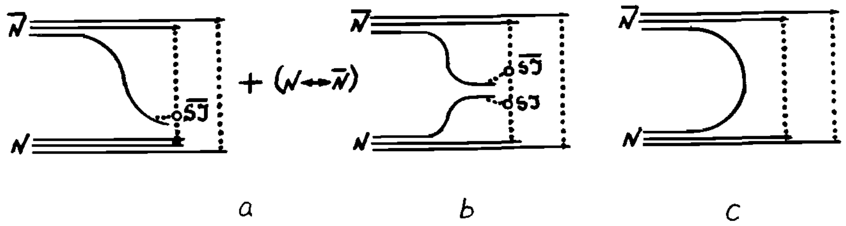

Fig. 4. Final state string configurations, corresponding to processes, shown in fig. 3.

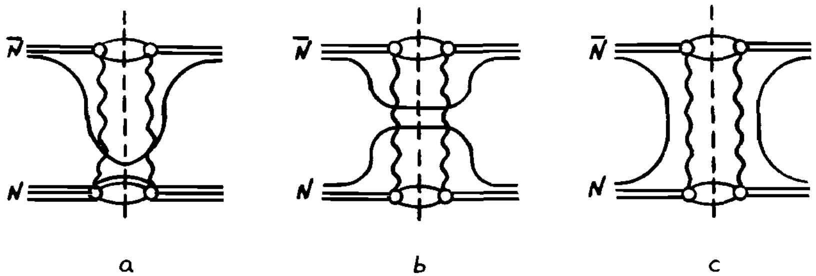

Fig. 5. The contributions to the $\bar{N}N$ elastic scattering amplitude, related by unitarity with the processes shown in fig. 3.

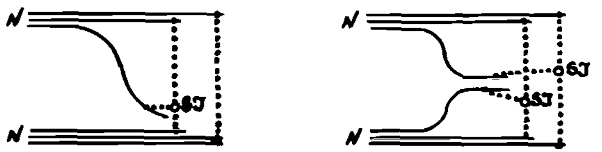

Fig. 6. Final state string configurations in those NN interaction channels, which compensate in $\Delta\sigma^{pp}_{tot}$ the fig. 3 a, b mechanisms contribution.

The concept of SJ is natural in the string model of nucleon which has a configuration of Y [17], where SJ couples all three strings together. It also arises when one constructs the nucleon creation operator in the gauge invariant form [15,16]. Unfortunately, the method of topological classification of the annihilation graphs is not supported by any calculation scheme. As a rule the magnitudes of the M^J_4, M^J_2 and M^J_0 Regge-trajectory intercepts are adopted [15,16,18] to be equal to -0.5, 0 and 0.5 respectively. This estimation has been deduced by Eilon and Harrari [19] from the following relations, obtained in the bootstrape multiperipheral model [20,21]:

$$\alpha_{M^J_4}(0) = 2\alpha_B(0) - \alpha_R(0) \quad , \qquad (1)$$

$$\alpha_{M^J_2}(0) = 2\alpha_B(0) - 2\alpha_R(0) + 1 \quad , \qquad (2)$$

$$\alpha_{M^J_0}(0) = 2\alpha_B(0) - 3\alpha_R(0) + 2 \quad , \qquad (3)$$

where the baryonic and mesonic Regge-intercepts have been fixed by the values $\alpha_B(0)=0$ and $\alpha_R(0)=0.5$. The expressions (1)-(3) can be also obtained by means of the Kaidalov's relation [22], following from the assumptions on s-channel factorization and on the SJ existence. The most reliable is the estimation of $\alpha_{M^J_4}(0)$, because the diquarks are maintained unharmed on the graph in fig. 1a, whereas they are destroyed on fig. 1b,c.

The paper [19] by Eilon and Harrari is a source of the widespread opinion [1,5,15,16,18,23,24] that the observed energy dependence (up to 12 GeV) of the $\bar{p}p$ annihilation cross section $\sigma^{\bar{p}p}_{ann} \propto$ $\propto s^{-1/2}$ is connected with the asymptotical contribution of the three-sheet graph in fig. 2 c.

In the model of $\bar{N}N$ annihilation by Rossi and Veneziano [15,16] the properties of the $\bar{q}q$-Reggeons strictly differ from those, which are expected in the planar Reggeon model. Indeed, the difference $\Delta\sigma^{pp}_{tot} = \sigma^{\bar{p}p}_{tot} - \sigma^{pp}_{tot}$ is written in the approach [15,16] as follows

$$\Delta\sigma^{pp}_{tot} = 2\sigma_\omega + \sigma_J \, , \qquad (4)$$

where σ_ω and σ_J are the contributions to $\Delta\sigma^{pp}_{tot}$ of ω-Reggeon and of all the $M^J_{4,2,0}$-Reggeons respectively. The experimentally observed relation $\Delta\sigma^{pp}_{tot} \approx \sigma^{\bar{p}p}_{ann}$ and formula (4) imply the strong suppression of ω-contribution: $2\sigma_\omega \ll \Delta\sigma^{pp}_{tot}$. On the contrary, the planar Reggeon model gives the relation [4] $\sigma_\omega \approx 9\sigma_\rho$, which leads to the estimation $2\sigma_\omega \approx (0.6 \div 0.8)\Delta\sigma^{pp}_{tot}$ by using the value of σ_ρ obtained from the analysis of total cross sections [25]. About the same value of σ_ρ follows from the factorization relation $\sigma^{NN}_\rho = (\sigma^{\pi N}_\rho)^2/\sigma^{\pi\pi}_\rho$, and from the estimation of $\sigma^{\pi\pi}_\rho$, obtained in the analysis [26] of the inelastic πN-scattering data. Thus, the planar structure of $\bar{q}q$-Reggeons, supposed in the QGSM, is inconsistent with the model of $\bar{N}N$-annihilation proposed in [15,16].

The suppression $\sigma_\omega \ll \sigma_J$ comes also into contradiction with the widely adopted explanation of the observed growth of pp inclusive cross section in the central rapidity region with the energy. It is connected in the QGSM [3] with a flow of the valence quark into the central rapidity region, i.e. with a contribution of a cilindrical Pomeron graph, containing one undeveloped sheet. Avoiding the unnatural assumption on the dominant probability of the $\overline{SJ}$-SJ annihilation in comparison with the $\bar{q}$-q one, we can find, that in the case of $\sigma_\omega \ll \sigma_J$, SJ is flowed into the central rapidity region much more frequently, than the valence quarks do. But if SJ acquires the rapidity $y \approx 0$, the final state has the three-sheet configuration in some rapidity region, and it results in a decreasing energy dependence of the inclusive cross section. Note, that the flow of the SJ through large rapidity gap has been considered earlier by Volkovistsky [24] in connection with the problem of the Eilon-Harrari's compensation [19] mechanism. But the author has not paid attention to the above difficulty with the explanation of the growth of the inclusive spectra.

One more puzzle arises both in the approach [15,16] (irrespective of the number of the graph in fig. 2 a,b,c, which dominates the annihilation at the intermediate energies) and in the models [18,27], where annihilation is connected by the unitarity with the cutted

Reggeons. The fact is: the annihilation is connected with the small impact parameter region /18/, but the difference of the $\bar{p}p$ and pp elastic scattering amplitudes is very peripheral in the impact parameter plane. This contradiction can be avoided, as it was proposed in /15,16/, by mixing of the $\bar{q}q$- and $\overline{SJ}$-SJ-Reggeons. Nevertheless no numerical estimation, supporting this proposition, has been carried out. Thus, the problem of $\bar{p}p$-annihilation mechanism demands further investigation.

Earlier, in papers /28-30/, the authors considered the asymptotical $\bar{p}p$ annihilation mechanism, connected with formation of the three-string in final state due to the colour decuplet gluonic exchange in the cross channel. It corresponds, in a sense of the unitarity, to the cut of so-called decameron exchange graph /30/. The estimation in the perturbative QCD (double-gluon exchange) provided the cross section of the decuplet-antidecuplet final state production of about $\sigma_{\{10\}} \approx 1 \div 2$ mb. Although the perturbative QCD is formally invalid for soft hadronic processes, some kind of expansion over small parameter can be carried out. Indeed the decuplet exchange is accompanied by the diquark destruction, i.e. transition from antitriplet to sixtet colour state. It leads to the dissapearance of $\sigma_{\{10\}}$, when the mean distance between any pair of quarks in the nucleon tends to zero. Thus, the results /28/ can be treated as a first order approximation in perturbative expansion over small parameter r_D^2 / r_p^2, where r_D^2 and r_p^2 are the mean squares of the diquark and proton radii. Although in the simplest case of oscillatory nucleon wave function (WF) $r_D^2 / r_p^2 = 3/4$, there are many evidences of admixture in the nucleon WF of a compact diquark with $r_D^2 / r_p^2 \approx 0.2$. These are the results of the deep inelastic scattering /31-34/, the investigation of the processes with large p_T hadron production /35-37/, the calculations in the model of instanton vacuum /38/, etc. These circumstances allow one to hope for the correct estimation of $\sigma^{\bar{p}p}_{ann}$ in the perturbative QCD. In any case, application of the perturbative theory methods to the annihilation has more ground, than in the case of total cross section, where the double gluon approximation gives, nevertheless, good results /12,39,40/.

The considered above asymptotical contribution to the annihilation cross section slightly decreases with the energy, mainly due to the possibility of $\bar{N}N$ sea-pairs production and the decameron Reggeization. Such small asymptotical tail cannot be picked out yet from the present experimental data (up to 12 GeV), which can be approximated /18,29,30/ by $\sigma^{\bar{p}p}_{ann} \approx C(S/S_0)^{-0.56}$, where C = 70 mb, S_0 = 1 GeV^2.

Nevertheless, the analysis /29,30/ of the experimental data on $\bar{p}p$ and pp multiplicity distribution up to ISR energies, picked out the decameron contribution to the high multiplicity events with the energy independent cross section $\sigma_{\{10\}}$ = 1.5 ± 0.1 mb. This value agrees very well with the perturbative QCD estimation. The observed decreasing energy dependence of $\sigma^{\bar{p}p}_{ann}$, as was proposed in /28,29/, can be connected with the diagrams in fig. 3, which contain transition of diquarks into sixtet colour states and there exists also the primordial valence $\bar{q}q$, $\bar{q}D$ or $\bar{D}q$ pairs with small ($\sim m_N$) invariant mass (fig. 3 a, b), or pair annihilates (fig. 3 b). The presence of the valence slowed down quark in the projectile nucleon WF leads to a correct energy-dependence $\sigma^{pp}_{ann} \propto S^{-1/2}$ for all the mechanisms, shown in fig. 3.

In the next section of present paper we formulate an approach, which combines methods of the perturbative QCD with the QGSM. We estimate in this way the contribution of the graphs in fig. 3 a,b to the annihilation cross section in section 3. The essential property of this mechanisms is their participation, in a sense of the unitarity, in the Pomeron, i.e. they satisfy the Eilon-Harrari's hypothesis. On the contrary, the process shown in fig. 3 c has no analogy with any channel in NN interaction. The corresponding cross section is estimated in section 4. This mechanism plays a role of nonplanar screening of the ordinary Reggeons. In other words, its contribution to $\Delta\sigma^{pp}_{tot}$ is negative. The sum of all the contributions of the graphs in fig. 3 a, b, c turn out to be in agreement with the experimental data. Some experimental evidences of the double - sheet topology of mechanisms, dominating annihilation at the energies about 10 GeV, are presented in section 5.

2. Perturbative QCD and soft hadronization

It is supposed usually in the QGSM /1-8/ that the admixture of valence gluons in the nucleon WF is negligibly small. This assumption is supported by a good agreement of the S-channel factorization relation /22/ with the experimental data. The former is based on the following space-time pattern of the Reggeon exchange. The annihilation of the valence $\bar{q}q$ pair is followed by formation of the universal colour string, which properties are independent of flavours of the annihilating $\bar{q}q$ pair. The hadronization of the colour strings is also supposed independent of the history of their creation and is carried out with unit probability.

Thus, one can estimate in the perturbative QCD the production cross section of such valence quarks configurations, which can decay then into pure mesonic final state. This cross section can be identified with the annihilation one (with correction for the branching of annihilation channel, of course). Such two-step approach is the result of the trouble with using the 1/N classification of annihilation mechanisms as it is done in the case of inelastic collision of mesons.

On the other hand, such separation into phases is dictated by characteristic properties of the $\bar{N}N$ annihilation. Only the destruction of one of the incident diquark and antidiquark at least makes it possible to have the annihilation. As the admixture of constituent gluons has been neglected, no primordial sixtet diquark can exist, but they should be prepared in a short time interval along with the $\bar{N}N$ interaction. It is clear, that the diquark destruction can be evaluated in one-gluon approximation, for the compact diquark at least. The interaction by means of Coulomb gluon exchange is instantaneous in the contrary to the soft hadronization process, which proceeds long time. This allows one to consider the two phases of the annihilation reaction (this separation is obviously invariant under the longitudinal Lorenz boosts).

For the sake of simplicity we consider the sixtet-antisixtet $D_{\{6\}} - \bar{D}_{\{\bar{6}\}}$ system as two triplet colour strings, as it is dictated by 1/N expansion in the QCD. This simplification is not principal. It is natural, that tension of the colour sixtet string is larger than the triplet one. In this case the successive decay of the sixtet string by means of $\bar{q}q$ pairs, tunneling from vacuum (we ignore the $\bar{D}D$ creation), leads to the same final hadronic state, as two triplet string do.

3. Contribution of the annihilation channels to the Pomeron

Let us consider the graphs shown in fig. 3 a, b. This graphs are connected by the unitarity with the Pomeron, because single gluon is only exchanged in t-channel. On the other hand, they give a contribution to the annihilation channels. Indeed, the graph in fig. 3 a has to do with the case, when one of the valence projectile antiquarks is situated in the nucleon's fragmentation region close to the diquark in the rapidity scale, and vice versa, the valence quark and antidiquark have close rapidities. In the gluon exchange between fast colour triplet antidiquark $\bar{D}_{\{3\}}$ and nucleon converts the former into antisixtet colour state $\bar{D}_{\{\bar{6}\}}$, the configuration of strings, shown in fig. 4 a, appears in the final state (dotted lines denote strings). We consider, as it was mentioned above,

the decay of the chromoelectric tube with $\{\bar{6}\}$ - $\{6\}$ on its ends, as independent hadronization of two triplet strings, according to the 1/N-method prescription. The position of the SJ in the string configuration M_4^J in fig. 4 a is determined by condition on smallness of the full string energy. This condition forbids the decay of the configuration M_4^J into states with antinucleon in the projectile $\bar{N}$ fragmentation region.

The graph in fig. 3 b corresponds to the case, when the valence $\bar{q}q$ pair has small relative rapidity interval. The valence colour triplet $\bar{D}_{\{3\}}$ and antitriplet $D_{\{\bar{3}\}}$ interact by means of gluon-exchange, and turn into $\bar{D}_{\{\bar{6}\}}$ and $D_{\{6\}}$ states respectively. The corresponding configuration of strings in the final state is shown in fig. 4 b. The configuration M_4^J also appears here.

In view of absence of a consistent theory, describing the processes of formation and decay of the strings, it is possible now to carry out only the crude estimation of the annihilation cross section, corresponding to the string configuration in fig. 4 a, b. It is natural to assume, that the probability of decay of this string configuration M_4^J only into mesons is close to unity (we neglect everywhere the vacuum $\bar{N}N$ pair tunneling production. It is suppressed at the energies below few tenth of GeV, and can be easily taken into account at higher energies). Indeed, after the long parts of the strings in the configuration M_4^J are broken into pieces, the following decay of M_4^J into $\bar{N}N$ can be forbidden by energy conservation, if the effective mass of each $\bar{q}D$, $q\bar{D}$ or $\bar{q}q$ pairs is small enough. However the short link of the string between $\bar{S}\bar{J}$ and SJ can be broken first. This suppresses the annihilation by factor 2 ÷ 3. On the other hand, the configuration M_4^J can also decay into pure mesonic state above the energy threshold. So, we choose the limit $M_{D\bar{q},q\bar{q}} \leq m_N$ in calculation of the annihilation cross section, corresponding to the graph in fig. 3 a, b .

Let us start with the graph in fig. 3 a. Its contribution to the annihilation cross section can be estimated as follows

$$\sigma_{ann}^{(3a)} \approx 2 P_{\bar{q}D}(M)\, \sigma_{\{6\}} \cdot \qquad (5)$$

Here $P_{\bar{q}D}(M)$ is a probability to find in the initial state a pair of valence $\bar{q}D$ with an invariant mass, which is smaller than M ; $\sigma_{\{6\}}$ is a cross section of $\bar{D}N$ interaction by means of one-gluon exchange, which results in transformation of $\bar{D}_{\{3\}}$ into antisixtet colour state $\bar{D}_{\{\bar{6}\}}$. If a nucleon is considered as qD system, we put on the supplementary condition: diquark should be remained in the triplet colour state. Factor 2 in (5) takes into account the possibility of primordial small mass $q\bar{D}$ pair, which correspond to the interchange $N \rightleftharpoons N$ in fig. 3 a.

The cross section $\sigma_{\{6\}}$ has in the one-gluon approximation the form

$$\sigma_{\{6\}} = \frac{\alpha_s^2}{4}\int\frac{d^2\vec{q}}{q^4}\sum_{\{6\},\{\bar{3}\}}\langle\bar{\mathcal{D}}_{\{3\}}|\sum_{i=1}^{2}(\hat{\lambda}_i^{\alpha})^T e^{i\vec{q}\vec{r}_i}|\bar{\mathcal{D}}_{\{\bar{6}\}}\rangle \times$$

$$\langle\bar{\mathcal{D}}_{\{\bar{6}\}}|\sum_{j=1}^{2}(\hat{\lambda}_j^{\beta})^T e^{-i\vec{q}\vec{r}_j}|\bar{\mathcal{D}}_{\{3\}}\rangle\langle N|\sum_{k=1}^{3}\hat{\lambda}_k^{\alpha}e^{-i\vec{q}\vec{r}_k}|q\mathcal{D}_{\{\bar{3}\}}\rangle \times \tag{6}$$

$$\langle q\mathcal{D}_{\{\bar{3}\}}|\sum_{\ell=1}^{3}\hat{\lambda}_\ell^{\beta}e^{i\vec{q}\vec{r}_\ell}|N\rangle .$$

Here α_s is the QCD coupling constant; $\hat{\lambda}^{\alpha}$ are the Gell-Mann's colour matrices.

The summing up over spin and isospin parts of WF of $\bar{D}_{\{\bar{6}\}}$ and $D_{\{\bar{3}\}}$ can be performed in (6) without any restrictions, imposed by the Pauli principle, because the expressions inside the brakets in (6) are symmetrical with respect to the permutation of quarks. The sum over diquark colour indices has the form

$$\sum_{\{6\}}\langle i,j|\{6\}\rangle\langle\{6\}|k,\ell\rangle = \frac{1}{2}(\delta_k^i\delta_\ell^j + \delta_\ell^i\delta_k^j),$$

$$\sum_{\{3\}}\langle i,j|\{3\}\rangle\langle\{3\}|k,\ell\rangle = \frac{1}{2}(\delta_k^i\delta_\ell^j - \delta_\ell^i\delta_k^j), \tag{7}$$

where i, j, k, l are the quark colour indices.

Using these formulae together with the relation

$$\langle i,j|\{3\}_k\rangle = \frac{1}{\sqrt{2}}\varepsilon_{ijk} , \tag{8}$$

we can transform expression (6) into the following one

$$\sigma_{\{6\}} = \frac{4\pi\alpha_s^2}{3}\int_0^\infty\frac{dq^2}{q^4}[1-F_2^D(q)]\left\{\frac{3}{2}[1-F_2^N(q)F_1^D(q)] + \right. \tag{9}$$

$$\left.\frac{1}{2}[F_2^D(q) - F_2^N(q)F_1^D(q)]\right\} .$$

Here $F_{1,2}^D(q)$ are the single-quark and double-quark form factors of diquark, which are defined as

$$F_1^D(q) = \langle\mathcal{D}_{\{3\}}|e^{i\vec{q}\vec{r}_1}|\mathcal{D}_{\{3\}}\rangle$$

$$F_2^D(q) = \langle\mathcal{D}_{\{3\}}|e^{i\vec{q}(\vec{r}_1-\vec{r}_2)}|\mathcal{D}_{\{3\}}\rangle . \tag{10}$$

The nucleon is considered here as a q-D system, so form factor $F_2^N(q)$ has a form

$$F_2^N(q) = \langle N|e^{i\vec{q}(\vec{r}_q-\vec{r}_D)}|N\rangle . \tag{11}$$

Expression (9) for $\sigma_{\{6\}}$ is infrared stable. It is the result of the fact that the long-waved gluon cannot resolve the inner structure of a diquark. For this reason we put effective gluon mass $M_g = 0$.

We use nonrelativistic oscillatory form of WF of diquark and of q-D relative motion in a nucleon to calculate $F_{1,2}^D$ and F_2^N :

$$F_i^D(q^2) = exp(-\alpha_i^D q^2), \tag{12}$$

where $\alpha_1^D = r_D^2/6$, $\alpha_2^D = 2r_D^2/3$;

$$F_2^N(q^2) = exp(-\alpha_2^N q^2), \tag{13}$$

where $\alpha_2^N = \frac{2}{3}(r_p^2 - r_D^2/3)$.

Here r_D^2 is a diquark mean radius squared; r_p is a nucleon charge radius, which is fixed by the value 0.8 F. To calculate $F_2^N(q^2)$ we supposed the dominance of the component with scalar diquark of mass $m_D \approx m_q$ /41/ in the nucleon WF. However it is inessential for the value of $\sigma_{\{6\}}$, which is insensitive to the form of the nucleon WF.

Expression (9) with form factors (12), (13) takes a form

$$\sigma_{\{6\}} = \frac{4\pi\alpha_s^2}{3}\left\{\frac{3}{2}J(\alpha_2^{\bar{D}}, \alpha_1^D + \alpha_2^N) + \right. \tag{14}$$

$$\left.\frac{1}{2}I(\alpha_2^D, \alpha_2^{\bar{D}}, \alpha_1^D + \alpha_2^N - \alpha_2^D)\right\} ,$$

where

$$\mathcal{J}(a,b) = a\ln\left(\frac{a+b}{a}\right) + b\ln\left(\frac{a+b}{b}\right)$$

$$I(a,b,c) = (a+b)\ln\left(\frac{a+b+c}{a+b}\right) + (a+c)\ln\left(\frac{a+b+c}{a+c}\right) - a\ln\left(\frac{a+b+c}{a}\right).$$

Parameter α_i^D differs from $\alpha_i^{\bar{D}}$ in (14), because we intend to take into account the distribution over diquark dimension in the nucleon WF.

To fix the coupling α_s we normalize the cross section of the process $NN \to N_{\{8\}}N_{\{8\}}$, calculated in the one-gluon exchange approximation

$$\sigma(NN \to N_{\{8\}}N_{\{8\}}) = 8\pi\alpha_s^2 \int \frac{dq^2}{q^4}\left[1 - \langle N| e^{i\vec{q}(\vec{r}_1-\vec{r}_2)}|N\rangle\right]^2, \tag{15}$$

to the experimental value of σ_{in}^{NN}. For the oscillatory nucleon WF $\langle N|\exp[i\vec{q}(\vec{r}_1-\vec{r}_2)]|N\rangle = \exp(-r_p^2 q^2/2)$, so

$$\alpha_s = \left(\frac{\sigma_{in}^{NN}}{8\pi\mathcal{J}\left(\frac{r_p^2}{2}, \frac{r_p^2}{2}\right)}\right)^{\frac{1}{2}}. \tag{16}$$

Putting σ_{in}^{NN} = 30 mb, one finds α_s= 0.52. More cumbersome formulae for $\sigma(NN \to N_{\{8\}}N_{\{8\}})$ in the case of asymmetrical WF with compact diquark, give close value of α_s (if r_p is fixed).

The calculation of $\sigma_{\{6\}}$ are performed with two types of the nucleon WF. In the first case (denoted by I) we put $r_D \approx 0.7$ F (it corresponds to symmetrical oscillatory nucleon WF). In the second case (II) we suppose the existence of about 50% admixture of the component with compact diquark /31,35/ with $r_D \approx 0.4$ F. These two possibilities result in the following values of $\sigma_{\{6\}}$

$$\sigma_{\{6\}} \approx \begin{cases} 0.31\cdot\sigma_{in}^{NN} \approx 9 \text{ mb} & \text{(I)} \\ 0.23\,\sigma_{in}^{NN} \approx 7 \text{ mb} & \text{(II)}. \end{cases} \tag{17}$$

The top and bottom values correspond everywhere to calculations in variants I and II respectively.

To calculate the annihilation cross section (5), one should now estimate the probability $P_{\bar{q}D}(M)$. We use the following parametrization /6/ for the N's antiquark distribution over the variable $x_+ = \frac{1}{2}(x + \sqrt{x^2 + x_\perp^2})$ $(x_\perp = 4m_\perp^2/s,\ m_\perp^2 = m_q^2 + \langle p_{q\perp}^2\rangle)$:

$$f_{\bar{q}}(x_+) = C\, x_+^{-\alpha_R(0)}(1-x_+)^{\beta-1}, \tag{18}$$

where $\beta = 1 + \alpha_R(0) - 2\alpha_N(0)$; $\alpha_R(0) = 0.5$; $\alpha_N(0) = -0.5$. The normalizing factor C is defined by the condition of the unit probability to find antiquark anywhere in the N's WF :

$$C^{-1} = \frac{\Gamma[1-\alpha_R(0)]\Gamma(\beta)}{\Gamma[\beta-\alpha_R(0)+1]} - \frac{1}{1-\alpha_R(0)}\left(\frac{m_q^2 + p_{q\perp}^2}{s\,|x_{\bar{q}}^{min}|}\right)^{1-\alpha_R(0)} \tag{19}$$

$x_{\bar{q}}^{min}$ is a minimum value of the Feynman variable of antiquark, which is thrown to the nucleon fragmentation region (we define $x_{\bar{N}}$ = 1). The configurations with leading diquark dominate in this region. So we put $x_{\bar{q}}^{min} \approx -0.5$.

Unfortunately, the correct description of the nucleon fragmentation region with antiquark thrown here, in particular, the correlation of momenta in the $\bar{q}qD$ system, seems to be impossible now. The formula (19) has, nevertheless, a weak dependence on $x_{\bar{q}}^{min}$ and reflects correctly to some extent the diminution of the phase space of incident quarks with the energy decreasing.

The distribution (18) can be rewritten for $x_+ \to 0$ in the rapidity scale in c.m.s. as

$$W_{\bar{q}}(y_{\bar{q}}) = C\left(\frac{m_\perp}{\sqrt{s}} e^y\right)^{1-\alpha_R(0)}. \tag{20}$$

Neglecting the q-D momentum correlation, one can write the probability $P_{\bar{q}D}(M)$ in the form

$$P_{\bar{q}D}(M) \approx \int_{y_{\bar{q}}^{min}+\delta(M)}^{y_{\bar{q}}^{min}} dy\, W_{\bar{q}}(y), \tag{21}$$

where

$$\delta(M) = \operatorname{arch}\left(\frac{M^2 - m_{q\perp}^2 - m_{D\perp}^2}{2m_{q\perp}m_{D\perp}}\right). \tag{22}$$

The minimum value of antiquark rapidity $y_{\bar{q}}^{min}$ is determined by the condition $(m_q+m_D)\exp|y_{\bar{q}}^{min}|/\sqrt{s} \approx 1$, which follows from the restriction on Feynman variable of $\bar{q}$D-system: $|x_{\bar{q}D}| \lesssim 1$ (we take into account in (21) the leading character of diquark in the nucleon fragmentation region, so $|y_D| > |y_{\bar{q}}|$).

Finally, we put $M = 1$ GeV and $m_D = m_q = 0.3$ GeV, in accordance with the results of nucleon Regge-trajectory fit /42/ ($m_D \approx 0.22$ GeV), and calculations /43/ in the instanton vacuum model. Then the expression (21) yields at the energy ~ 10 GeV

$$P_{\bar{q}D}(M = 1\,GeV) \approx 1.7\,(s/s_0)^{-1/2} \tag{23}$$

$$s_0 = 1\ GeV^2.$$

Now, we can join the expressions (5), (17) and (23) to obtain an estimation of the contribution to σ_{ann} of the graph in fig.3a

$$\sigma_{ann}^{(3a)}(\bar{N}N) \approx \begin{cases} \sqrt{\frac{s_0}{s}}\,30 \text{ mb} & \text{(I)} \\ \sqrt{\frac{s_0}{s}}\,24 \text{ mb} & \text{(II)}. \end{cases} \tag{24}$$

We should turn further to the graph in fig. 3 b. The corresponding contribution to $\sigma_{ann}^{\bar{P}P}$ is written in the form, analogous to (5):

$$\sigma_{ann}^{(3b)} \approx P_{\bar{q}q}(M)\,\sigma_{\{\bar{6},6\}}\,. \tag{25}$$

Here $P_{\bar{q}q}(M)$ is a probability of configurations, where mass of the projectile $\bar{q}q$ pair is smaller than M. $\sigma_{\{6,6\}}$ is the cross section of the process $\bar{D}_{\{3\}}D_{\{\bar{3}\}} \to \bar{D}_{\{\bar{6}\}}D_{\{6\}}$ averaged over the triplet colour indices of $\bar{D}_{\{3\}}$ and $D_{\{\bar{3}\}}$ states:

$$\sigma_{\{\bar{6},6\}} = \frac{1}{9}\sum_{n,m=1}^{3}\sigma_{\{\bar{6},6\}}^{n,m}\,, \tag{26}$$

where

$$\sigma_{\{\bar{6},6\}}^{n,m} = \frac{\alpha_s^2}{4}\int\frac{d^2\vec{q}}{q^4}\sum_{\{\bar{6}\},\{6\}}\langle\mathcal{D}_{\{3\}_n}|\sum_{i=1}^{2}\lambda_i^{\alpha}e^{i\vec{q}\vec{\tau}_i}|\mathcal{D}_{\{6\}}\rangle \times$$
$$\langle\mathcal{D}_{\{6\}}|\sum_{j=1}^{2}\lambda_j^{\beta}e^{-i\vec{q}\vec{\tau}_j}|\mathcal{D}_{\{3\}_n}\rangle\langle\bar{\mathcal{D}}_{\{3\}_m}|\sum_{k=1}^{2}(\lambda_k^{\alpha})^T e^{i\vec{q}\vec{\tau}_k}|\bar{\mathcal{D}}_{\{\bar{6}\}}\rangle \times$$
$$\langle\bar{\mathcal{D}}_{\{\bar{6}\}}|\sum_{l=1}^{2}(\lambda_l^{\beta})^T e^{-i\vec{q}\vec{\tau}_l}|\bar{\mathcal{D}}_{\{3\}_m}\rangle\,. \tag{27}$$

The summing up over space and spin-isospin parts of the WF of $D_{\{6\}}$ and $\bar{D}_{\{\bar{6}\}}$ can be performed as in expression (6) without any restriction imposed by the Pauli principle. Taking into account the fact, that the colour WF of $\{3\}$ is antisymmetrical, but WF of $\{6\}$ is symmetrical, one can rewrite (27) in the form

$$\sigma_{\{\bar{6},6\}}^{n,m} = \frac{\pi\alpha_s^2}{4}C_{nm}\int\frac{d^2\vec{q}}{q^4}\left[1 - F_2^{D}(q)\right]\left[1 - F_2^{\bar{D}}(q)\right]. \tag{28}$$

The factors C_{nm} can be found, using (7), (8)

$$C_{nm} = 5 + 9\,\delta_{nm}\,. \tag{29}$$

Expression (28) is infrared stable as well as $\sigma_{\{6\}}$.

Using the formulae (12), (26), (28) and (29), we obtain

$$\sigma_{\{\bar{6},6\}} = \frac{\sigma_{in}^{NN}}{4}\,\frac{J(\frac{2}{3}\tau_D^2, \frac{2}{3}\tau_{\bar{D}}^2)}{J(\frac{1}{2}\tau_P^2, \frac{1}{2}\tau_P^2)}\,, \tag{30}$$

which is numerically equal to

$$\sigma_{\{\bar{6},6\}} \approx \begin{cases} 8 \text{ mb} & \text{(I)} \\ 5 \text{ mb} & \text{(II)}. \end{cases} \tag{31}$$

The factor $P_{\bar{q}q}(M)$ in (25) is estimated by means of the formula

$$P_{\bar{q}q}(M) \approx \int_{|y_{q,\bar{q}}|<Y} W_{\bar{q}}(y_{\bar{q}})W_q(y_q)\,\theta(\varepsilon - |y_q - y_{\bar{q}}|)\,dy_{\bar{q}}\,dy_q\,, \tag{32}$$

where $\varepsilon = 2\,\mathrm{arch}(M/2m_{q_\perp})$; $W_{\bar{q}}(y_{\bar{q}})$ is defined in (20); $W_q(y_q) = W_{\bar{q}}(-y_q)$; $Y = \ln(x_0\frac{\sqrt{s}}{m_{q_\perp}})$, $x_0 \approx 1/3$. The last factor arises naturally, when the smooth distribution (18) is changed to (20) in the nucleon fragmentation region.

The integral (32) is taken by means of changing the integration variables by $(y_{\bar{q}} \pm y_q)$. The restriction $|y_{\bar{q},q}| < Y$ is taken into account approximately by diminution of limits of integration over $(y_{\bar{q}} + y_q)/2$ in a value of $\langle|\Delta y|\rangle$, which is the mean value of the rapidity interval between $\bar{q}$ and q.

$$\langle|\Delta y|\rangle = \frac{\int_{-\varepsilon}^{\varepsilon} d\Delta y\cdot\Delta y\,\exp\left([1-\alpha_R(0)]\Delta y\right)}{\int_{-\varepsilon}^{\varepsilon}\exp\left([1-\alpha_R(0)]\Delta y\right)d\Delta y}\,. \tag{33}$$

One obtains in this way

$$P_{\bar{q}q}(M) = 4C\left(\frac{M^2-4m^2_{q_\perp}}{s}\right)^{\frac{1}{2}}\left\{\ln\left(\frac{2x_0\sqrt{s}}{M+\sqrt{M^2-4m^2_{q_\perp}}}\right)+\left(\frac{M-2m_{q_\perp}}{M+2m_{q_\perp}}\right)^{\frac{1}{2}}\right\}. \tag{34}$$

Finally, using (25), (31) and (34), we find the contribution of the graph in fig. 3 b to the annihilation cross section

$$\sigma^{(3b)}_{ann}(\bar{N}N) \approx \begin{cases} \sqrt{\frac{S_0}{S}}\,28\text{ mb} & \text{(I)} \\ \sqrt{\frac{S_0}{S}}\,18\text{ mb} & \text{(II)} \end{cases} \tag{35}$$

Thus, the sum of contributions (24) and (35) explains considerable part of the observed annihilation cross section in the energy range $E \sim 10$ GeV.

Contribution of the graphs in fig. 3 a, b to the total cross section shown in fig. 5 a, b, is a part of the Pomeron, as was noted in the beginning of this section, i.e. it satisfies the hypothesis by Eilon and Harrari [19]. The corresponding part of the Pomeron contribution to the $\bar{N}N$ total cross section is connected by the unitarity with processes shown in fig. 6 a, b. Those channels of $\bar{N}N$ interaction, which compensate the annihilation in the difference $\sigma^{\bar{p}p}_{tot} - \sigma^{pp}_{tot}$, contain a small effective mass $\bar{N}N$ pair in the fragmentation regions (fig. 6 a) or in central rapidity region (fig. 6 b).

4. Nonplanar annihilation corrections to Reggeons

The primordial valence $\bar{q}q$ pair in fig. 3 c annihilates, in contrast to fig. 3 b. Pair of spectators $\bar{D}D$ turns to colour sixtet-antisixtet system after the gluon exchange and is hadronized as two independent $\bar{q}q$ strings (fig. 4 c). The contribution of this mechanism to the annihilation cross section is connected with the cut of the diagram shown in fig. 5 c, where only the colour sixtet $\bar{D}D$ intermediate states are taken into account (as before the sea $\bar{N}N$ pairs are neglected).

There is no consistent theory of $\bar{q}q$ Reggeons, so we consider the diagram in fig. 5 c under two different assumptions on the dynamics of Reggeon exchange. The exchanged quarks are supposed in

variant A to emit and absorb soft gluons and this process is described by some planar diagram [44]. On the contrary, the colour index of the quark line is not changed in variant B, no gluons are emitted. The quark exchange in this case picks out some rare configuration in the nucleon WF [41].

Variant A. In the leading 1/N approximation one can neglect the change of the diquark triplet colour indices, after the double gluon exchange in fig. 5 c. As a result, the scattering amplitude, corresponding to the graph in fig. 5 c, can be written in the eikonal approximation as

$$\Delta_{\{\bar{6},6\}}T_R(\vec{b}) = -\Gamma_{\{\bar{6},6\}}(\vec{b})\,T^{\circ}_R(\vec{b}). \tag{36}$$

Here $\vec{b}$ is an impact parameter; $\Gamma_{\{\bar{6},6\}}(b)$ is a profile function of the double-gluon contribution to the diquark elastic scattering with $\{\bar{6}\}$ - $\{6\}$ intermediate state. The notation $\Delta_{\{\bar{6},6\}}T_R$ emphasizes the fact, that contribution (36) is a correction to the bare Reggeon amplitude $T^{\circ}_R(b)$. The expression (36) takes into account the fact that the centres of mass of the diquark with $x \approx 1$ and of nucleon coincide in the impact parameter plane.

Let us parametrize the functions $T^{\circ}_R(b)$ and $\Gamma_{\{\bar{6},6\}}(b)$ in the Gaussian form

$$\Gamma_{\{\bar{6},6\}}(\vec{b}) = \frac{\sigma_{\{\bar{6},6\}}}{8\pi\lambda_{\{\bar{6},6\}}}\exp\left(-\frac{\vec{b}^2}{4\lambda_{\{\bar{6},6\}}}\right). \tag{37}$$

$$T^{\circ}_R(\vec{b}) = T^{\circ}_R(0)\exp\left(-\frac{\vec{b}^2}{4\lambda_R}\right). \tag{38}$$

The parameter $\lambda_{\{\bar{6},6\}}$ in (37) can be evaluated as $\lambda_{\{\bar{6},6\}} \approx r^2_D/3$. The exact value of $\lambda_{\{\bar{6},6\}}$ is unessential, because of smallness of the diquark radius in comparison with the interaction radius of the Reggeon amplitude $T^{\circ}_R(\vec{b})$.

Now we can estimate the contribution of the graph in fig. 5 c to the imaginary part of the NN forward scattering amplitude

$$\Delta_{\{\bar{6},6\}}T_R(q=0) = -K^R_{\{\bar{6},6\}}T^{\circ}_R(q=0),$$

where

$$K^{R}_{\{\bar 6,6\}} = \frac{\sigma_{\{\bar 6,6\}}}{8\pi(\lambda_R+\lambda_{\{\bar 6,6\}})} \qquad (39)$$

The phenomenological analysis /45/ of data at energies 10 ÷ 20 GeV provides the values $\lambda_f \approx 5\ (\mathrm{GeV/c})^{-2}$, $\lambda_\omega \approx 11\ (\mathrm{GeV/c})^{-2}$. Then (39) and (31) give

$$K^{f}_{\{\bar 6,6\}} \approx \begin{cases} 0.1 & (\mathrm{I}) \\ 0.066 & (\mathrm{II}) \end{cases} \qquad (40)$$

$$K^{\omega}_{\{\bar 6,6\}} \approx \begin{cases} 0.065 & (\mathrm{I}) \\ 0.045 & (\mathrm{II}) \end{cases} \qquad (41)$$

To find the corresponding contribution to the annihilation cross section, one can use the AGK cutting rules /46/ (normalization is $\sigma_R = \mathrm{Im}T_R/s$):

$$\sigma^{(3c)}_{ann}(\bar N N) = 2K^{f}_{\{\bar 6,6\}}\sigma^{o}_{f} + 2K^{\omega}_{\{\bar 6,6\}}\sigma^{o}_{\omega}\,. \qquad (42)$$

Now, we are faced with a necessity of estimation of the bare Regge cross sections σ^{o}_{f} and σ^{o}_{ω} in order to find $\sigma^{(3c)}_{ann}$ by using (40), (41).

Let us start with determination of σ^{o}_{ω}. It is known only an effective ω-contribution from the experiment :

$$\Delta\sigma^{pp}_{tot} = 2\sigma^{eff}_{\omega} = 2(\sigma^{o}_{\omega} - \sigma^{abs}_{\omega})\,, \qquad (43)$$

where σ^{abs}_{ω} contains except $\Delta_{\{\bar 6,6\}}\sigma_\omega$ other corrections, considered below.

The contamination of the sea makes it possible to produce $\{\bar 6\}-\{\bar 3\}$ or $\{3\}-\{6\}$ intermediate colour states in addition to $\{\bar 6\}-\{6\}$ state, considered above. These intermediate states are related by the unitarity with inelastic $\bar N N$ interactions, where valence nucleon or antinucleon are thrown to the central rapidity region.

The contribution of the graph in fig. 5 c with the $\{\bar 6\}-\{\bar 3\}$ and $\{3\}-\{6\}$ intermediate states to a screening correction to the bare Regge amplitude T^{o}_{ω}, can be estimated also by means of formula (39), where factor $\sigma_{\{\bar 6,6\}}$ is changed by $2\sigma_{\{3,6\}}$. The last one is a cross section of the process $\bar D_{\{3\}}D_{\{3\}} \to \bar D_{\{\bar 6\}}D_{\{\bar 3\}}$, averaged

over colour indices of initial particles. It is given by the following expression

$$\sigma_{\{3,6\}} = \frac{\alpha_s^2}{36}\sum_{n,m}\int\frac{d^2\vec q}{(q^2+M_g^2)^2}\sum_{\{6\},\{3\}}\langle \mathcal{D}_{\{\bar 3\}_n}|\sum_{i=1}^{2}\lambda_i^\alpha e^{i\vec q\vec\tau_i}|\mathcal{D}_{\{6\}}\rangle \times$$
$$\langle \mathcal{D}_{\{6\}}|\sum_{j=1}^{2}\lambda_j^\beta e^{i\vec q\vec\tau_j}|\mathcal{D}_{\{\bar 3\}_n}\rangle\langle \bar{\mathcal{D}}_{\{3\}_m}|\sum_{k=1}^{2}(\lambda_k^\alpha)^T e^{-i\vec q\vec\tau_k}|\bar{\mathcal{D}}_{\{3'\}}\rangle \times \qquad (44)$$
$$\langle \bar{\mathcal{D}}_{\{3'\}}|\sum_{\ell=1}^{2}(\lambda_\ell^\beta)^T e^{-i\vec q\vec\tau_\ell}|\bar{\mathcal{D}}_{\{3\}_m}\rangle\,.$$

We are forced to introduce the effective gluon mass M_g here, because $\sigma_{\{3,6\}}$ is not infrared stable on the contrary to $\sigma_{\{\bar 6,6\}}$.

By means of relations (3) and (4), expression (44) is reduced to

$$\sigma_{\{3,6\}} = \frac{2\pi\alpha_s^2}{3}\int_0^\infty\frac{dq^2}{(q^2+M_g^2)^2}\left[1-F_2^{\mathcal{D}}(q)\right]\left[1-F_2^{\bar{\mathcal{D}}}(q)\right] \qquad (45)$$

If one puts here $F_2^D = F_2^{\bar D}$ for the sake of simplisity and uses parametrization (12), he finds

$$\sigma_{\{3,6\}} = \frac{8\pi\alpha_s^2}{9}\tau_D^2\exp\left(\tfrac{4}{3}\tau_D^2M_g^2\right)Ei\left(-\tfrac{4}{3}\tau_D^2M_g^2\right)\,. \qquad (46)$$

The value of α_s is determined, as usual, by means of normalization to $\sigma^{NN}_{in} \approx 30$ mb. We put $M_g \approx m_\pi$, fix $\alpha_s \approx 0.7$ and find from (46) and (39)

$$K^{\omega}_{\{3,6\}} = \begin{cases} 0.14 & (\mathrm{I}) \\ 0.1 & (\mathrm{II})\,. \end{cases} \qquad (47)$$

Let us gather all the screening correction (41), (47) and ordinary $\omega\mathbb{P}$-cuts (without distruction of triplet diquarks), which have $K_{\omega\mathbb{P}} \approx 0.25$. Then the bare ω-Reggeon cross section is appeared to be connected with the experimentally measured one, due to (43), as

$$\sigma^{o}_{\omega}/\Delta\sigma^{pp}_{tot} \approx \begin{cases} 0.9 & (\mathrm{I}) \\ 0.83 & (\mathrm{II})\,. \end{cases} \qquad (48)$$

The connection between the bare f and ω Reggeon contributions depends on a degree of the exchange degeneracy violation. The fit of the experimental data in the quasieikonal model /45/ has given $\sigma_f/\sigma_\omega \approx 2.7$. In this case the contribution of the mechanism, shown in fig. 3 c, to the annihilation cross section in variant A is found from (40) - (42) to be equal to

$$\left[\sigma_{ann}^{(3c)}(\bar{N}N)/\Delta\sigma_{tot}^{pp}\right]_A \approx \begin{cases} 0.6 & (I) \\ 0.37 & (II) \end{cases} . \qquad (49)$$

Variant B. No accompanying gluon is allowed now. Hence, only the intermediate state $\{\bar{6}\}$ - $\{6\}$ is possible on the graph in fig. 5 c. Once more difference of variant B, is supplementary colour factor 7/4, which should be introduced into right hand side of expression (39). One can be convinced of this, considering the formula (29) for the factors C_{nm}. The procedure analogous to variant A leads to the following estimation

$$\left[\sigma_{ann}^{(3c)}(\bar{N}N)/\Delta\sigma_{tot}(pp)\right]_B \approx \begin{cases} 0.95 & (I) \\ 0.55 & (II) \end{cases} . \qquad (50)$$

It is worthy to emphasize, that the results (49)-(50) should be considered as an order of magnitude estimation only. They depend strongly on the exchange degeneracy violation degree. The phenomenological approach [47], based on Pomeron structure in perturbative QCD [48,49], describes well the present-day accelerator data on σ_{tot}^{pp}, the cosmic ray data and the recent $S\bar{p}pS$ result on the elastic scattering phase [50]. It gives a weak violation of the exchange degeneracy [50]: $\sigma_f/\sigma_\omega \approx 1.2$. In this case the numbers in the right-hand sides of (49), (50) should be about two times smaller.

5. Some consequences of the topological structure of the annihilation graphs

1. Some peculiarities of the annihilation channels have been found [51] at the momentum p_{lab} = 32 GeV/c: as the multiplicity of the particles produced is enlarged, the geometry of the annihilation events becomes more planar. The enlarged multiplicity corresponds in the framework of the QGSM to the more uniform division of the energy between the strings. The transverse momenta of the quarks settled on the string ends in fig. 4 can be shown to be determined mainly by the gluon momentum in fig. 3. It leads to the correlation between the directions of the jets in the forward and backward hemispheres. The larger is the multiplicity of the event, the more pronounced is the four-jet structure of the event in the c.m.s., the larger is its planarity.

On the contrary, the mechanism corresponding to three-sheet topology causes the decrease of the planarity with the multiplicity growth for the same reasons. Thus, three-sheet dominance contradicts the data [51].

It is worth noting that the multiplicity in the analysed events was changed in data [51] from 6 to 12 particles. This is just the region of dominance of the two-sheet mechanisms as it follows from the fit of multiplicity distributions, carried out in our papers [29,30]. The large multiplicity events are determined in part by the decameron contribution, corresponding to the three-sheet topology. There are some other multi-sheet contributions (absorbtion corrections caused by multi-Pomeron rescattering, for instance), located in the high multiplicity region. Thus, we expect the planarity to fall at larger than [51] multiplicities.

The Fermi-motion of valence quarks inside the incident hadrons smears out partly the multiplicity-planarity correlation. For this reason, it is better to investigate the geometrical structure of the events separately for the forward and backward hemispheres in the c.m.s. but not for the whole totality of hadrons as it was done in [51]. One can expect the V-like structure of the events, caused by the diquark distruction.

2. The experimentally observed up to 12 GeV ratio of the annihilation mean multiplicity to the total one $\langle n \rangle_{ann}/\langle n \rangle_{pp} \approx 3/2$ is considered [1,8,15,16,24,27] generally as an important support of the dominant contribution to the annihilation of the three-sheet mechanism. However, we show here the fulfillment of this relation at intermediate energies for the two-sheet topology too. Let us use the formulae for the mean multiplicity of particles, produced in the decay of the $\bar{q}$-q and q-D strings of the mass $\sqrt{S}$, which were obtained in paper [6] by the analysis of the data on e^+e^--annihilation and deep-inelastic scattering.

$$\begin{aligned} \langle n \rangle_{\bar{q}-q} &= 2.1 + 0.85 \ln (S/S_0) \\ \langle n \rangle_{q-D} &= 0.8 + 0.9 \ln (S/S_0) . \end{aligned} \qquad (51)$$

Here S_0 = 1 $(GeV)^2$. The large value of $\langle n \rangle_{\bar{q}-q}$ in comparison with $\langle n \rangle_{q-D}$ results in the increase of the annihilation multiplicity, which is indeed found approximately 1.5 times larger than the total one.

Thus, the enhanced annihilation multiplicity cannot be considered as an argument in favour of the three-sheet topology dominance.

It is a trivial consequence of the fact, that the energy is not high enough.

6. Conclusions

Let us review shortly the main result of the present investigation.

1. There are two principially different in nature, but comparable in a value types of contributions to the $\bar{N}N$ annihilation cross section at intermediate energies.

The first type of processes is shown in fig. 3 a, b and 4 a, b. It is connected by unitarity with a cut of the Pomeron contribution to the $\bar{N}N$ elastic scattering amplitude, shown in fig. 5 a, b. These mechanisms satisfy approximately the compensation hypothesis by Eilon and Harrari. The corresponding contribution to the NN interaction cross section, which compensates the annihilation in the difference $\Delta\sigma^{pp}_{tot} = \sigma^{\bar{p}p}_{tot} - \sigma^{pp}_{tot}$, is shown in fig. 6 a, b. A pair of nucleons, unconnected with the sea $\bar{N}N$ pair production, appears with a small effective mass in the final state of this processes.

The second type of annihilation processes is shown in fig. 3 c, 4 c. It has no matter with the Pomeron amplitude, but plays a role of nonplanar correction to the Reggeon amplitude. This contribution to the $\bar{N}N$ annihilation cross section violates the hypothesis by Eilon-Harrari. However it does not enlarge $\Delta\sigma^{pp}_{tot}$, but on the contrary suppresses it, because this correction to the ω-Reggeon is of shadowing type. Thus, the more central is the annihilation in the impact parameter plane, the more peripheral is the difference of the $\bar{p}p$ and pp elastic amplitudes. This conclusion removes the well known contradiction.

2. All mechanisms, shown in fig. 3, are connected with the slowing down of valence quark (antiquark). It naturally explains the experimentally observed energy dependence of the annihilation cross section $\sim s^{-1/2}$.

3. The sum of estimated contributions of the mechanisms in fig. 3 a, b, c to the annihilation cross section is comparable with the value, given by experiment. The well known relation $\sigma^{pp}_{ann} \approx \Delta\sigma^{pp}_{tot}$ is seemed to be accidental. Indeed, as all the mechanisms in fig. 3 a, b, c are connected with diquark destruction, then their contribution to σ^{NN}_{ann} should dissapear, when diquark radius tends to zero. In the same time, $\Delta\sigma^{pp}_{tot}$ is almost independent on diquark radius.

4. All the considered mechanisms have two-sheet topology. The exchange of semi-hard gluon results in correlation between asimuthal angles of the planes, which are formed by V-type events in forward and backward hemispheres. The larger is multiplicity the more is planarity of events.

Small diquark radius and connected with it large transverse momentum of gluon lead to considerable sea-gull effect, i.e. fast growth of the mean transverse momenta of particles $\langle p_T(x)\rangle$ versus Feynman variable. It is observed experimentally indeed /18/.

5. One of the by-products of present consideration is a nontrivial mechanism of flowing the baryon quantum number through a large rapidity gap by means of single-quark exchange. It will be considered in detail in a separate paper.

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Расщепление дикварков и предасимптотические механизмы аннигиляции антипротонов

Рассмотрен ряд предасимптотических механизмов $\bar{N}N$-аннигиляции. Общим для них является наличие медленного кварка /антикварка/ в волновых функциях налетающих адронов и расщепление одного или обоих дикварков путем перевода их из триплетного по цвету состояния в сикстетное. Некоторые аннигиляционные механизмы являются частью померонного вклада в сечение $\bar{N}N$-взаимодействия, другие играют роль непланарной экранировки реджеонных амплитуд. Суммарный вклад механизмов в сечение аннигиляции хорошо соответствует данным как по величине, так и по энергетической зависимости.

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Diquark Destruction and Preasymptotical Mechanisms of Antiproton Annihilation

A few preasymptotical mechanisms of $\bar{N}N$ annihilation are considered. Their common feature is the presence of slowed down valence quark /antiquark/ in the wave functions of projectiles and the destruction of one or both diquarks by means of transformation from triplet to sixtet colour state. One class of annihilation mechanisms is a part of the Pomeron amplitude, but another one plays a role of nonplanar shadowing to Reggeon amplitude. The sum of all the mechanism contributions is in good agreement with the experimentally measured annihilation cross section both in the value and in the energy dependence.

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

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