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COLOUR DYNAMICS IN LARGE p_T HADRON PRODUCTION ON NUCLEI

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1. As is widely accepted, in quantum chromodynamics the main contribution to the inelastic cross section of hadrons is given by a colour exchange. The produced colour objects fly away with large relative momenta and a colour string (tube) is stretched between them. The breaking of the string by spontaneous $q\bar{q}$ pair production from the vacuum in the colour field leads to multiple hadron production [2,3].

However, for considerations below it will be important only that the fast end of the string is retarded by a constant force κ. For colour triplets at the ends of the string we have $\kappa \approx (2\pi\alpha'_R)^{-1} \approx 1$ GeV/fm.

2. Before hard collision on a given nucleon of the nucleus the incident hadron may undergo inelastic interaction on the surface of the nucleus, i.e., exchange colour. The coloured hadron loses its momentum, $\Delta p = \kappa\Delta z$, on a distance Δz. Hence the hard scattering takes place at a smaller incident energy and as $x_T \to 1$ this is only allowed on the surface of the nucleus.

3. Let a quark with momentum k, produced in a hard collision, fragments into a hadron with momentum βk at a distance ℓ_f from the interaction point. The retardation of the quark by the colour string gives for the leading hadron ($\beta \geq 1/2$) that $\ell_f \approx \frac{k}{\kappa}(1-\beta)$. Subsequent colour exchanges of the quark on the nucleons of the nucleus are not essential, since they "rotate" it only in the colour space, i.e., the quark is not absorbed by the nucleus over distance ℓ_f.

As $x_T \to 1$ we have also $\beta \to 1$, hence $\ell_f \to 0$. Consequently, the hadron in this case is formed immediately after the hard scattering of the quark and it can be absorbed by the nucleus. From this and section 2 we conclude that the cross section for producing a hadron with large p_T (in the case of single hard scattering) depends on A as $\sim A^{1/3}$ when $x_T \to 1$.

4. To pick out single hard scattering inside the nucleus, in experiments [1,5] the authors studied the production of symmetric, large-p_T hadron pairs produced at angle 90° in the c.m.s. of the incident hadron and a nucleon of the target nucleus. Denote by a_1 and a_2 the momentum fractions of quarks from the incident proton and the nucleon of the nucleus undergoing hard collision. The scattered quarks fragment into mesons carrying momentum fractions β_1 and β_2 of the quark momenta. If one neglects the transversal motion of quarks in the incident hadrons and transversal momenta of hadrons in final jets, then one has $a_1 \approx a_2 = a$, $\beta_1 \approx \beta_2 = \beta$, $x_T \approx a\beta$.

The invariant cross section $E_1E_2 d\sigma/d^3p_1d^3p_2 \equiv \sigma_{inv}(x_T)$ for production of a large-p_T hadron pair in NN collision is given by

$$\sigma_{inv}(x_T) = C\int da_1 da_2 d\beta_1 d\beta_2 F(a_1)F(a_2)D(\beta_1)D(\beta_2) \times \times \sigma_{qq}(a_1a_2 s)\delta(x_T - a_1\beta_1)\delta(x_T - a_2\beta_2). \quad (1)$$

Here $x_T = 2p_T/\sqrt{s}$; $s = 2m^2 + 2mE$, where E is the energy of the incident proton, $F(a)$ is the momentum distribution of quarks within the nucleons; we take $F(a) \propto (1-a)^3/\sqrt{a}$.

Similarly, for the quark fragmentation functions $D(\beta) \propto (1-\beta)^2/\beta$. The elastic quark-quark scattering cross section has the form $\sigma_{qq}(Q) \propto Q^{-4}$ in the one-gluon-exchange approximation; C is the normalization constant.

In the case of nuclear target we single out two contributions to the symmetric hadron-pair-production cross section: when the incident hadron has no inelastic interactions before the hard collision ($\sigma^{(1)}_{inv}$) and when it exchanges colour on one or more nucleons of the nucleus ($\sigma^{(2)}_{inv}$).

$$\sigma^{(1)}_{inv}(x_T) = C\int d^2b \int_{-\infty}^{+\infty} dz\,\rho(b,z)\exp[-\sigma^{NN}_{in}T(b,-\infty,z)] \times \times \int_{x_T}^{1} da\,F^2(a)D^2\left(\frac{x_T}{a}\right)a^{-2}\sigma_{qq}(a^2 s)\exp[-(\sigma^{h_1N}_{in} + \sigma^{h_2N}_{in})T(b,z+\ell_f,\infty)]. \quad (2)$$

$$\sigma^{(2)}_{inv}(x_T) = C\int d^2b \int_{-\infty}^{+\infty} dy\,\sigma^{NN}_{in}\rho(b,y)\exp[-\sigma^{NN}_{in}T(b,-\infty,y)]\int_{y}^{\infty} dz\,\rho(b,z) \times \times \int_{\tilde{x}_T}^{1} da\,F^2(a)D^2(\tilde{x}_T/a)a^{-2}\sigma_{qq}(a^2\tilde{s})\exp[-(\sigma^{h_1N}_{in} + \sigma^{h_2N}_{in})T(b,z+\tilde{\ell}_f,\infty)]. \quad (3)$$

In these expressions we have done the integrations over a_2, β_1, β_2. We have also introduced the following notation: $\rho(b,z)$ is the nuclear density depending on impact parameter b and longitudinal coordinate z;

$$T(b,z_1,z_2) \equiv \int_{z_1}^{z_2} dz\,\rho(b,z); \quad \tilde{s} = 2m^2 + 2m\tilde{E}; \quad \tilde{E} = E - \kappa(z-y);$$

$$\tilde{x}_T = 2p_T/\sqrt{\tilde{s}}; \quad \ell_f = E(a - x_T)/(2\kappa); \quad \tilde{\ell}_f = \tilde{E}(a - \tilde{x}_T)/(2\kappa).$$

By using (2) and (3) one can determine the A-dependence of the cross sections. The result of calculations for

$\alpha_{12}(x_T) = d\ln[\sigma^{(1)}_{inv}(x_T) + \sigma^{(2)}_{inv}(x_T)]/d\ln A$

at A = 100, E = 70 GeV, κ = 3 GeV/fm is shown in fig.1 by the dotted line. The nuclear density was taken in the Woods-Saxon form. As it is seen, by increasing x_T, the exponent $\alpha_{12}(x_T)$ drops until ~1/3.

5. At moderate p_T values it is possible that the two hadrons are formed independently, being in a symmetric configuration only by chance. The corresponding contribution to the invariant cross section is:

$$\sigma^{(3)}_{inv}(x_T) \approx E_1 \frac{d^3\sigma}{d^3p_1} E_2 \frac{d^3\sigma}{d^3p_2} / \sigma^{NA}_{in} . \qquad (4)$$

The A dependence of this contribution is given by the exponent $\alpha_3(x_T) = d\ln[\sigma^{(3)}_{inv}(x_T)]/d\ln A \approx \alpha_{h_1}(x_T) + \alpha_{h_2}(x_T) - 2/3$.

Here $\alpha_h(x_T)$ is the corresponding exponent for the inclusive production of hadron h on nuclei. The values of $\alpha_h(x_T)$ for different hadrons were taken from results of ref.[6] at E = 70 GeV.

6. The cross section for hadron pair production in NN collisions is approximated by the expression[6]:

$$\sigma^{NN}_{inv}(x_T) = h e^{-Hp_T} + r e^{-Rp_T}. \qquad (5)$$

The first term corresponds to hard scattering, the second to a random pair production. This interpretation is supported by the behaviour of the correlation function[7], which increases sharply at $p_T \geq 0.7$ GeV/c. The exponent $\alpha_{eff}(x)$ characterizing the A dependence of $\sigma^{NA}_{inv}(x_T)$ can be estimated by the formula:

$$\alpha_{eff}(x_T) = \frac{\alpha_{12} + \alpha_3 A^{\alpha_3-\alpha_{12}} e^{(H-R)p_T} r/h}{1 + A^{\alpha_3-\alpha_{12}} e^{(H-R)p_T} r/h} . \qquad (6)$$

In our calculations we took for simplicity the same values of r/h and H-R for all types of particles and fixed them in accordance with experimental data[7] as r/h = 0.006 and H-R = = -6 $(GeV/c)^{-1}$. The results of calculations for $\pi^+\pi^+$ and π^+K^+ pairs are shown in figs.1,2 by solid lines. The sharp change of α_{eff} at $p_T \approx 1$ GeV/c is due to the fact that the hard scattering becomes dominant.

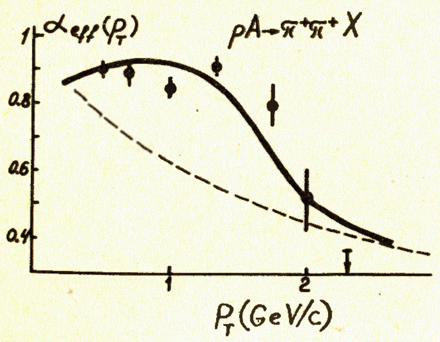

Fig.1

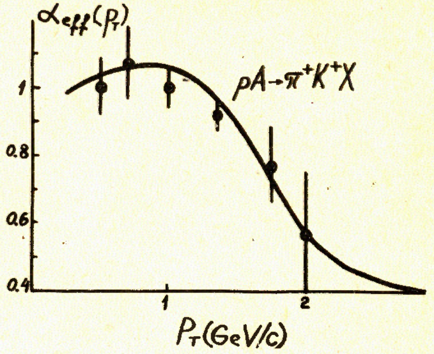

Fig.2

7. The experimental data[6,7] show that large-p_T protons are produced with a larger cross section than pions. This fact indicates that the large-p_T proton is produced not by fragmentation of a quark, but due to the collective scattering of the three-quark system with large momentum transfer. It is clear that such a process singles out in the initial proton configurations possessing a small size $R \approx 1/p_T$. The absorption cross section for such states is small[8,9], $\tilde{\sigma}_{in}/\sigma_{in} \approx m^2_\pi/p^2_T$. Hence, in the case of NN pair production the incident and outgoing nucleons are not absorbed in the nucleus. Consequently, in (6) one has to take $\alpha_{12}(x_T) = 1$. The result of corresponding calculation is shown in fig.3.

In the case of πN or KN pairs the outgoing nucleon can be originated from the beam or target. In the first case in (2) and (3) one has to put $\sigma^{NN}_{in} = 0$ which leads to $\alpha' \approx 1$. In the second case $\sigma^{NN}_{in} = 0$ only for the outgoing nucleon, and the corresponding exponent α'' is approximately equal to that calculated for the $\pi\pi$ and πK pairs. For $\alpha_{eff}(x_T)$ in this case one has to put into (6) the value $\alpha_{12} = (\alpha' A^{\alpha'} + \alpha'' A^{\alpha''})/(A^{\alpha'} + A^{\alpha''})$. The results of calculation are shown in figs.4 and 5.

8. Evidently, the effects of retardation by colour forces become less important at a larger incident energy. For example, at E = 400 GeV $\alpha_{12}(x_T)$ starts to decrease only at $x_T \geq 0.7$, i.e., at $p_T \geq 10$ GeV/c. The available data[12] are for much smaller values of p_T, and hence, $\alpha_{eff} \approx 1$.

9. In the case of pion beam one can make the following predictions. If in the produced pair of hadrons one has a pion of the same charge as in the incident beam, then $\alpha_{eff} \approx 1$. In the opposite case $\alpha_{eff}(x_T)$ should fall off with increasing p_T like in the reaction $pA \to \pi\pi X$.

10. The results of experiment[1] turn out to be quite fruitful in understanding colour dynamics of hadronic interactions. In the framework of the presented interpretation they have demonstrated that colour objects are retarded by colour fields and also that compressed hadronic configurations are not absorbed by nuclear matter.

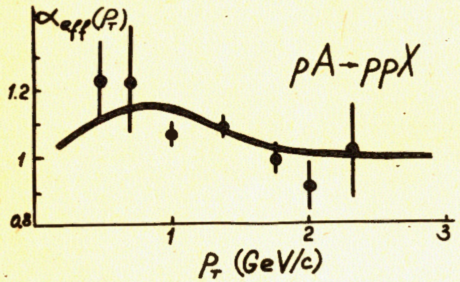

Fig.3

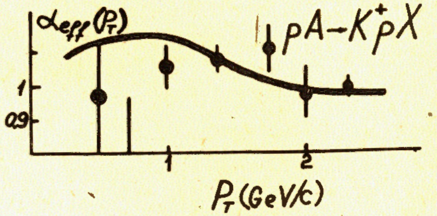

Fig.4

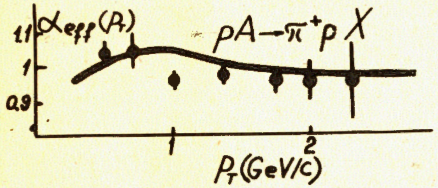

Fig.5

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Динамика цвета в образовании адронов с большими p_T на ядрах

Торможение цветных объектов, распространяющихся внутри ядра, приводит к сильному экранированию жестких процессов вблизи кинематической границы. Это объясняет слабую A-зависимость сечения образования на ядрах пар мезонов с большими p_T, обнаруженную в работе[1]. Малость сечения поглощения в ядре сжатых адронных конфигураций объясняет линейную A-зависимость образования pp-пар.

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