The considered dynamic effect can be used as a sensitive method of detecting acoustic nuclear resonance, while the static effect is suitable for a determination of the quadrupole moments of excited states of nuclei.

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ENERGY DEPENDENCE OF INCLUSIVE CROSS SECTION OF PARTICLE PRODUCTION IN THE PIONIZATION REGION

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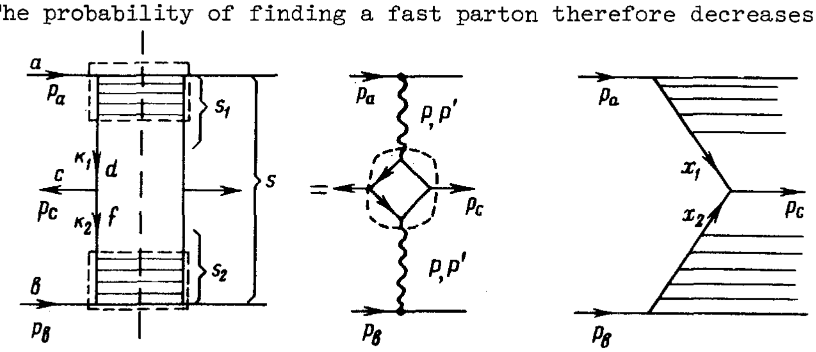
We consider the dependence of the inclusive cross section $f(p_c, s) = E_c d\sigma/d^3p_c$ of the production of pions that are slow in the c.m.s. on the colliding-hadron energy $\sqrt{s}$, and show that the experimentally observed growth of the cross section may be connected with the influence of phase space on the registered particles (Botke's numerical calculations [1] have shown that the influence of phase space extends to very large $s \sim 2000$ GeV^2). The point is that whereas scaling is well satisfied in the fragmentation region, even starting with accelerator energies [2] ($f(p_c, s) = f(x, p_\perp, s) = f(x, p_\perp, \infty)$ at $s > 20$ GeV^2), in the pionization region ($x = p_{\parallel c}/p_{\parallel max} = 0$) the value of $f(p_c, s)$ and the cross section for particle production at 90°, $d\sigma/d\Omega|_{90° cms} = \int f(p_{\parallel c} = 0, p_{\perp c}, s)\, p_\perp^2 dp_\perp/E_c$ increase with increasing s, up to the energies of the ISR colliding beams (thus, at $\sqrt{s} = 30$ GeV, $d\sigma/d\Omega|_{90°}$ is approximately twice as large as at $\sqrt{s} = 6.8$ GeV) [3, 4]. This effect is usually attributed [5] to the contribution of non-vacuum poles (P', ω) with $\alpha(0) = 1/2$ (see Fig. 1a), which fades out very slowly with increasing s, like $s_1^{-1/2} \sim s^{1/4}$. ($s_1 \sim \sqrt{s}$ in the case of 90° particle emission.) But then, to describe the growth of the cross section with energy, the vertex PP' (in a dashed box in Fig. 1a) must be chosen negative. This sign is quite strange if the two-reggeon PP' vertex is expressed in terms of the reggeon residues (as is done in Fig. 1a) and if it is recalled that in all known two-particle reactions the contribution of poles with $\alpha(0) = 1/2$ is positive (or is equal to zero if the channel is exotic [6].

On the other hand, since at accelerator energies ($s \sim 40$ GeV^2) the masses of shower particles with momenta larger than $p_c - \sqrt{s_1}$ and less than $p_c - \sqrt{s_2}$ are quite small ($s_1 \sim s_2 \sim 3$ GeV^2), it is necessary to take phase-space effects into account; these are easiest to explain with the parton model [7].

Particle "c" in Fig. 1b is formed upon collision of two partons. One of the hadrons, "a," carries a fraction x_1 of the momentum ($k_1 = x_1 p_a$), while the second "b" carries x_2. If parton 1 is fast enough ($x_1 \to 1$), then the remaining partons from "a" receive a small fraction of the energy, $\sim 1 - x_1$. The probability of finding a fast parton therefore decreases with increasing



a b c

Fig. 1

x_1. For example, in deeply inelastic ep scattering we have $\nu W_2(x) \to (1 - x)^3$ at $x = 1/\omega \to 1$. In our case, for emission of particle "c" we need $x_1 \sim x_2 \sim \sqrt{m_c^2 + p_{\perp c}^2}/\sqrt{s} + m_{\perp c}/\sqrt{s}$, and we can expect

$$f(0, p_\perp, s) \approx (1 - x_1)^3(1 - x_2)^3 f(0, p_\perp, \infty) \approx \left(1 - \frac{3m_{\perp c}}{\sqrt{s}}\right)\left(1 - \frac{3m_{\perp c}}{\sqrt{s}}\right) f(0, p_\perp, \infty) \qquad (1)$$

We shall now explain the cause of this dependence with the diagrams. The diagram of Fig. 1a corresponds to the integral

$$f(p_c\, s) = \frac{E_c\, d\sigma}{d^3 p_c} = g^2 \int \frac{d^4 k_1}{(2\pi)^4 s}\, 2\,\mathrm{Im} A_1 \,\mathrm{Im} A_2 \,|\,\eta\,(k_1^2)\,|^2 |\,\eta\,(k_2^2)\,|^2, \qquad (2)$$

where $\eta(k^2)$ is the propagator of particles "d" and "f," Im A_1 is the imaginary part of the forward scattering amplitude of hadron "d" by target "a," and g is the coupling constant.

The exact form of the propagator is immaterial. We see only that owing to the fast decrease of $\eta(k^2)$ with k^2 the integral with respect to k^2 converges at values $k^2 \sim 1/2R^2$. For concreteness, we can assume

$$\eta\,(k^2) = \exp\,(R^2 k^2), \quad (k^2 < 0). \qquad (3)$$

Using the optical theorem, we express the imaginary part of the scattering amplitude in terms of the total cross section for the interaction of particles a and d or f and b

$$\mathrm{Im}\, A_1 = (s_1 - M_n^2)\,\sigma_0\,(1 + 1/\sqrt{s_1}). \qquad (4)$$

The factor $s_1 - M_n^2$ in (4) (instead of the usual Im $A_1 = s_1\sigma(s_1)$) is a reflection of the threshold behavior of the amplitude. Im A_1 should vanish at $s_1 \leq (m_a + m_d)^2$, and $(1 + 1/\sqrt{s_1})$ describes the change of the cross section at high energies. Owing to the non-vacuum poles, $\sigma(s_1)$ approaches its asymptotic value from above.

It must be emphasized that expression (4) is not an expansion in Regge poles. This is simply a convenient approximation of the amplitude of meson-nucleon scattering (e.g., πN or γN).

It will be convenient in what follows to introduce the Sudakov variables [8]

$$p_a' = (0, 1, 2, 3) = \frac{\sqrt{s}}{2}(1, 0, 0, 1), \qquad p_b' = \frac{\sqrt{s}}{2}(1, 0, 0, -1),$$

$$p_c = \alpha_c p_a' + \beta_c p_b' + p_{\perp c}, \quad k_i = \alpha_i p_a' - \beta_i p_b' + k_{i\perp}.$$

Formula (2) then takes the form

$$f(p_c, s) = g^2 \int \frac{d\alpha_2 d\beta_1 d^2 k_{1\perp}}{(2\pi)^4}\, \sigma_0^2\, (s_1 - M_\Pi^2)(s_2 - M_\Pi^2)\left(1 + \frac{1}{\sqrt{s_1}}\right) \times$$
$$\times \left(1 + \frac{1}{\sqrt{s_2}}\right) |\,\eta\,(k_1^2)\,\eta\,(k_2^2)\,|^2\,, \qquad (5)$$

where all the corrections of order $1/\sqrt{s}$ are connected with the factor $(s_{1,2} - M_\pi^2)$. Indeed, $s_1 - M_\pi^2 = \beta_1(1 - \alpha_1)s - |k_{1\perp}|^2 + m_a^2 - M_\pi^2$. The important quantities in the integral (5) are $\beta_1 \approx 1/(2R^2 s\alpha_1)$, $\alpha_2 \approx 1/(2R^2 s\beta_2)$, $k_{1,2\perp}^2 \approx 1/2R^2$ (since $-k_1^2 = \alpha_1\beta_1 s + |k_{1\perp}^2|$ and $k_1^2 \approx 1/2R^2$ (see (3). Hence

$$s_1 - M_\Pi^2 = \beta_1 s\,[\,1 - \alpha_1 - \alpha_1(M_\Pi^2 - m_a^2)\,2R^2 - \alpha_1 2R^2/2R^2] \approx \beta_1 s(1 - 3{,}2\,\alpha_1)$$

$$s_2 - M_\Pi^2 \approx (1 - 3{,}2\beta_2)\,\alpha_2 s\,.$$

(Here $R^2 = 2\ \text{GeV}^{-2}$, $m_a^2 = m_N^2 = 0.9\ \text{GeV}^2$, $M_\pi^2 = (m_N + m_\pi)^2 = 1.2\ \text{GeV}^2$). We proceed similarly with the terms of order $s^{-1/4}$

$$\left(1 + \frac{1}{\sqrt{s_1}}\right)\left(1 + \frac{1}{\sqrt{s_2}}\right) \approx (1 + \sqrt{2R^2\beta_2})(1 + \sqrt{2R^2\alpha_1}) = (1 + 2\sqrt{\beta_2})(1 + 2\sqrt{\alpha_1}).$$

Finally, we estimate α_1 and β_1. The momentum conservation laws $\alpha_1 = \alpha_c + \alpha_2 \simeq \alpha_c + 1/(2R^2 s\beta_2)$ and $\beta_2 = \beta_c + \beta_1 \simeq \beta_c + 1/(2R^2 s\alpha_1)$ lead to the quantities

$$\alpha_1 = \alpha_c(1 + \delta), \quad \beta_2 = \beta_c(1 + \delta), \quad \delta = \sqrt{\frac{1}{4} + \frac{1}{2R^2 m_{\perp c}^2}} - \frac{1}{2} > 0. \tag{6}$$

As a result we obtain the formula

$$f(p_c, s) = (1 - 3{,}2\alpha_1)(1 - 3{,}2\beta_2)(1 + 2\sqrt{\beta_2})(1 + 2\sqrt{\alpha_1})\left[f(p_c, \infty) + 0\left(\frac{1}{s}\right)\right], \tag{7}$$

where the remaining integral

$$I = g^2 \int \frac{d\alpha_2 d\beta_1}{(2\pi)^4} d^2 k_{1\perp} \beta_1 \alpha_2 s^2 \sigma_0^2 |\eta(k_1^2)\eta(k_2^2)|^2 = f(p_c, \infty) + 0\left(\frac{1}{s}\right)$$

already contains corrections only of order 1/s and less.

In the derivation of (7) we specially did not put $\alpha_c = \beta_c$ (which is valid in the case of 90° emission), and left α_c and β_c arbitrary, because formula (7) describes not only the dependence of the inclusive cross section on the energy, but also the distribution with respect to the longitudinal rapidities of the particle "c" in the pionization region. The latter statement becomes obvious if we express α_c and β_c in terms of the rapidity of the hadron "c" in the c.m.s.

$$y_c = \frac{1}{2}\ln\frac{E_c + p_{\parallel c}}{E_c - p_{\parallel c}} = \frac{1}{2}\ln\frac{\alpha}{\beta}, \text{ i.e. } \alpha_c = \frac{m_{\perp c}}{\sqrt{s}} e^{y_c}, \quad \beta_c = \frac{m_{\perp c}}{\sqrt{s}} e^{-y_c}. \tag{8}$$

The energy dependence of $f(0, p_\perp, s)$ and the distribution with respect to the rapidity at various s, calculated with the aid of (7), (6), and (8), are shown in Figs. 2 and 3. The small humps (at $s \simeq 1000\ \text{GeV}^2$ in Fig. 2 and at $s = \infty$ in Fig. 3) are connected with the factor $(1 + 1/\sqrt{s_1})$ in (4), which takes into account the contribution of the non-vacuum poles (P', ω). However, owing to the threshold effects $(1 - 3.2\beta_2)$ the corrections of order $1/\sqrt{s_2}$ become noticeable only at very large s. At smaller s the humps overlap and form a plateau that is higher (by 40 - 50%) than the asymptotic one.

The parameters used in the calculation ($R^2 = 2\ \text{GeV}^{-2}$) were taken from [9], where a multiperipheral model with constant total cross section was constructed. We recognize in addition that an appreciable fraction of the pions is produced as a result of the decay of resonances (ρ mesons), and the longitudinal velocity y_π of the pion differs somewhat from y_ρ of the resonance. We therefore integrated formula (7) for these pions in the interval $y_\pi - 1 < y_\rho < y_\pi + 1$.

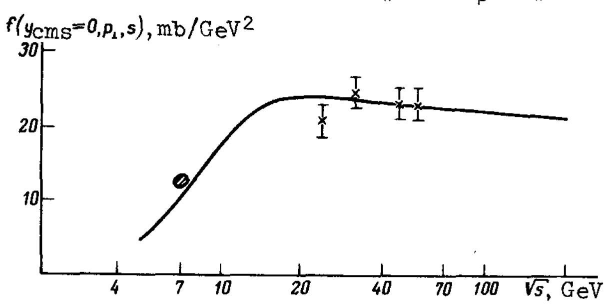


Fig. 2. Inclusive pp → π⁻ + X cross section vs energy ($p_\perp = 0.3$ GeV, $p_{\parallel cms} = 0$). The solid curve was calculated from formula (7). The circles and crosses represent the experimental data at accelerator [4] and ISR [3] energies, respectively.

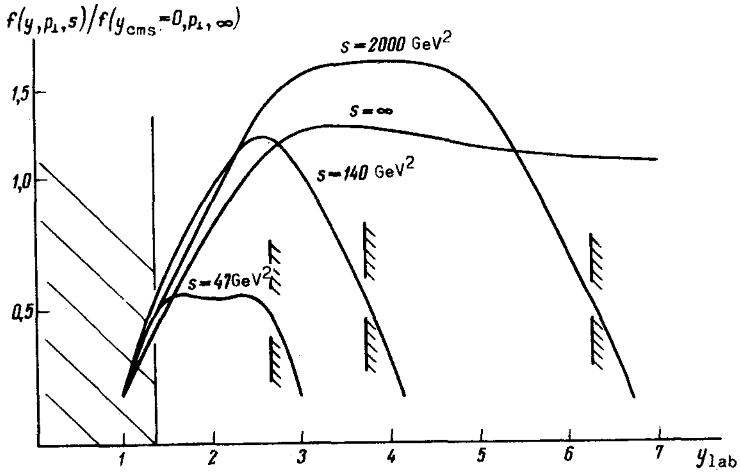


Fig. 3. Distribution of pions produced in the reaction $pp \to \pi^- X$ with respect to the longitudinal velocities ($p_\perp$ = 0.4 GeV). The solid curves were calculated from (7). The shaded region on the left is the region of fragmentation (three-reggeon limit), where formula (7) no longer holds, since β_2 is large.

The fraction of pions produced via resonances (= 2/3) and the average transverse resonance ass ($m_{\perp\rho}$ = 1 GeV) were also taken from [9].

We note finally an interesting qualitative effect observed by the Saclay-Strasbourg group [3]. With increasing $p_\perp$ (at $p_\perp$ > 500 MeV ≥ $\langle p_\perp \rangle$), the inclusive cross section begins to grow faster with increasing energy, owing to the increase of α_c and β_c in (7) with increasing $m_{\perp c} = [m_c^2 + p_{\perp c}^2]^{1/2}$.

We recall in conclusion that formula (7) does not claim to describe the experimental data exactly. We only wished to show that phase-space effects come into play even at sufficiently large s and can explain the role of the inclusive cross section in the pionization region. A detailed discussion of the energy dependence of $f(p_c, s)$ in a model of this type, with allowance for the decay of the resonances, will be published in the journal Yadernaya Fizika.

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POSSIBLE METHODS OF PLASTICITY CHANGE IN A SUPERCONDUCTING TRANSITION

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Dynamic mechanisms of weakening in N-S transitions, due to the decrease of electron dragging of dislocations, and quasi-static mechanisms due to the inhomogeneity of the dislocation structure, are discussed.

The transition of a superconductor from the normal (N) state to the superconducting (S) state at a temperature T below the critical temperature is accompanied by a jumplike increase of its plasticity (weakening [1 - 3]. This effect is observed under different deformation conditions, but always only during the stage of developed plastic deformation, and its connection with the presence of mobile dislocations in the crystal is subject to no doubt. It