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ЕФИ-12(72)

Yu.F.PIROGOV, V.A.KHOZE, N.L.TER-ISAAKIAN

ON POSSIBLE ELECTROMAGNETIC DOMINANCE NEAR THE KINEMATIC BOUNDARY OF HADRON SPECTRA

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Ю.Ф.ПИРОГОВ, В.А.ХОЗЕ, Н.Л.ТЕР-ИСААКЯН

О ВОЗМОЖНОСТИ ЭЛЕКТРОМАГНИТНОЙ ДОМИНАНТНОСТИ ВБЛИЗИ КИНЕМАТИЧЕСКОЙ ГРАНИЦЫ АДРОННЫХ СПЕКТРОВ

Показано, что при асимптотических энергиях в узкой области фазового объема вблизи кинематической границы в процессах рождения адронных пучков с относительно небольшими массами с ростом масс пучков и с уменьшением передаваемых импульсов начинают доминировать электромагнитные эффекты.

Указывается на необходимость учета этих эффектов при изучении адронных процессов, описываемых трехреджеонными формулами и их модификациями. Обсуждается также ситуация при больших, но фиксированных энергиях.

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Yu.F.PIROGOV, V.A.KHOZE, N.L.TER-ISAAKYAN

ON POSSIBLE ELECTROMAGNETIC DOMINANCE NEAR THE KINEMATIC BOUNDARY OF HADRON SPECTRA

It is shown that at asymptotic energies in the processes of hadron jet production with relatively low invariant masses, the electromagnetic effects begin to dominate in a narrow region of the phase space near the kinematic boundary as the beam masses increase and the momenta transfers decrease. The necessity to take into account these effects in studying the hadron processes described by the triple-Reggeon formula and their modifications is noted. The situation at higher but fixed energies is also described.

Yerevan Physical Institute

Yerevan 1972

It is of interest for the inelastic hadron processes theory the investigation of the asymptotic distribution near the phase-space boundary, since in this region, using the quasi-two-body phenomenology, the problem can be reduced to the determination of some unknown parameters, i.e., triple-Reggeon constants. For the particle beam production processes allowing vacuum exchange, as shown in Fig.1, in the limit $-t_i \lesssim m^2$; $s \gg M_1^2, M_2^2, M_{12}^2 \gg m^2$ the following triple-Reggeon formulae are valid [1,2]:

$$\frac{d\sigma^{(a)}_{3R}}{dt\,dM^2} = \frac{r_a(t)}{s^2}\left(\frac{s}{M^2}\right)^{2\alpha(t)}\left(M^2\right)^{\tilde\alpha(0)} \underset{t\to 0}{\approx} \frac{r_a(0)}{\left(M^2\right)^{2-\tilde\alpha(0)}} \qquad (1a)$$

$$\frac{d\sigma^{(b)}_{3R}}{dt\,dM_1^2\,dM_2^2} = \frac{r_b^2(t)}{s^2}\left(\frac{s}{M_1^2 M_2^2}\right)^{2\alpha(t)}\left(M_1^2 M_2^2\right)^{\tilde\alpha(0)} \underset{t\to 0}{\approx} \frac{r_b^2(0)}{\left(M_1^2 M_2^2\right)^{2-\tilde\alpha(0)}} \qquad (1b)$$

$$\frac{d\sigma^{(c)}_{3R}}{dt_1\,dt_2\,dM^2} = \frac{r_c^2(t_1,t_2)}{s^2}\int\frac{ds_1}{s_1}\left(\frac{s_1}{M^2}\right)^{2\alpha(t_1)}\left(\frac{s_2}{M^2}\right)^{2\alpha(t_2)}\left(M^2\right)^{\tilde\alpha(0)} \underset{t\to 0}{\approx}$$

$$\approx \frac{r_c^2(0,0)\ln\eta}{\left(M^2\right)^{2-\tilde\alpha(0)}}; \qquad \eta = \frac{s\sqrt{t_1 t_2}}{M^2 m_1 m_2} > 1 \qquad (1c)$$

In (1a-1c) are kept only the terms which do not vanish at $s \to \infty$ and correspond to the Pomeron exchange, $\tilde\alpha$ is a certain effective trajectory determining the behaviour of the total cross sections $\sigma^{(P)}_{tot}$ and $\sigma^{(PP)}_{tot}$, while $r_i(t_i)$ are the effective triple-Reggeon constants involving residues, signatures and kinematic factors.

In a series of papers [3,4], it has been shown that the triple-Pomeron constant is small ($\sim t^{1/2}$ at $t \to 0$) and, therefore, the contribution

of the triple-Pomeron mechanism (PP, P) to (Ia) is $\frac{d\sigma_{st}}{dt\, dM^2} \approx \frac{t\varphi'}{M^2}$.
According to the language of two-component duality, this means that the background production asymptotically vanishes and only the resonances or equivalently the trajectory P' ($\alpha(0) = \frac{1}{2}$) dual to the latter, give contribution to the processes under consideration. However, the experimental data [5] apparently show a more rapid decrease with increasing M [4]. Let us note that this decrease can be explaned by the 2π cut, corresponding to $\alpha(0) = -1$, which effectively takes into account the contribution to the matrix element of the processes $NN \to NN\pi$ of one-pion exchange [6] essential in the mass region $M \sim 1.2 \div 1.4$ GeV, and rapidly vanishes as M increases. In this connection we note that at $t \to 0$ with increasing M^2 the contribution of the electromagnetic effects becomes essential (even dominant).

Namely, the expressions for the cross sections of the electromagnetic processes, described by the diagrams of Fig.I, after substituting P for γ-quanta at $t \lesssim m_\pi^2$ in the region under consideration have the form:

$$\frac{d\sigma_{em}^{(a)}}{dt\, dM^2} = \frac{\alpha}{\pi} \frac{\sigma_T^{\gamma P}(M^2)}{|t|\, M^2}\left(1 - \frac{t_{min}}{t}\right); \quad |t_{min}| = \frac{m_N^2 M^4}{S^2} \tag{2a}$$

$$\frac{d\sigma_{em}^{(b)}}{dt\, dM_1^2\, dM_2^2} = \frac{\sigma_T^{\gamma P}(M_1^2)\, \sigma_T^{\gamma P}(M_2^2)}{8\pi^3\, M_1^2\, M_2^2}\left[1 + \left(1 - \frac{t_{min}}{t}\right)\right]; \quad |t_{min}| = \frac{M_1^2 M_2^2}{S} \tag{2b}$$

$$\frac{d\sigma_{em}^{(c)}}{dt_1\, dt_2\, dM^2} = \frac{\alpha^2 \sigma_T^{\gamma\gamma}(M^2)}{\pi^2 M^2 t_1 t_2}\left[(1 + \eta^{-2}) \ln \eta - (1 - \eta^{-2})\right] \tag{2c}$$

Here $\sigma_T^{\gamma p} \approx 110$ microbarn; $\sigma^{\gamma\gamma}(M^2) = \frac{(\sigma_T^{\gamma p})^2}{\sigma_T^{pp}} \approx 0.3$ microbarn. The results (2a - 2c) are obtained in the same way as those of [7]. The comparison of (Ia) and (2a) gives:

$$\frac{d\sigma_{em}^{(a)}}{d\sigma_{st}^{(a)}} = C_I \frac{(M^2)^{1-\tilde{\alpha}(0)}}{|t|}\left(1 - \frac{t_{min}}{t}\right) \qquad (3)$$

(here and later on M^2, S and t are given in $(\mathrm{Gev})^2$)

For the processes $pp \to pX$, $\pi p \to \pi X$ [4,5] one can obtain $C_p = (5 \pm 2.5)\, 10^{-5}$ and $C_\pi \approx 2C_p$ assuming $\tilde{\alpha}(0) = \frac{1}{2}$.

Really, however, the region, where the electromagnetic effects become appreciable, is essentially wider due to the interference between the hadron and electromagnetic mechanisms.

The contribution of interference, for instance, for the process (Ia) at high S has the form:

$$\frac{d\sigma_{int}^{(a)}}{dt\, dM^2} \approx \frac{\alpha F(t)}{S^2}\left(\frac{S}{M^2}\right)^{\alpha_p(t)} \frac{S\sqrt{1-\frac{t_{min}}{t}}}{M^2\sqrt{-t}}\, \mathrm{Re}\left[\eta(t,S,M^2)\right] \mathrm{Im}\, A_{\gamma p \to P_p}(M^2,t,0)\Big] \approx$$

$$\approx \frac{\alpha F(0)}{(M^2)^{1-\tilde{\alpha}(0)}} \left[\frac{\sqrt{1-\frac{t_{min}}{t}}}{\sqrt{-t}} \quad \mathrm{Re}\, \eta(t,S,M^2)\right]_{t\to 0} \qquad (4)$$

where $\eta(t,S,M^2)$ takes into account the contributions of the cuts, normal trajectories, the real part of the Pomeron, and the phase shift between strong and electromagnetic amplitudes, which are proportional to $\frac{1}{\ln^2 S/M^2}$, $\frac{(M^2)^{3/2-\tilde{\alpha}(0)}}{\sqrt{S}}$, t and $\alpha \ln t$, respectively, when $t \to 0$. Therefore, the contribution of interference can be neglected in the region of dominance.

Analogous situation occurs for the process of Fig.Ib. The corrections to the process of Fig.Ic are connected with the interference between P and γ exchanges, and are determined by the five-Reggeon diagrams with even number P

and have the following order of magnitude:

$$\frac{d\sigma}{dt_1\, dt_2\, dM^2} \underset{t_1 \sim t_2 \sim t \to 0}{\simeq} \frac{\alpha^2}{t\,(M^2)^{2-\bar{\alpha}(0)}} \qquad (5)$$

The whole foregoing consideration was of asymptotic nature assuming $S \to \infty$. However, at energies really achievable in the nearest future, it is necessary in addition to take into account the contributions of the normal trajectories. In this case, in view of the absence of interference between the vacuum and normal exchanges (P,R,P) and (PR,R) [8], and neglecting the term (RR,R), we have for the process (Ia) at $\bar{\alpha}(0) = \frac{1}{2}$ the following formula

$$\frac{d\sigma^{(a)}_{str}}{dt\, dM^2} \approx \left[\frac{\gamma_{PPR}(0)}{M^3} + \frac{\gamma_{RRP}(0)}{S}\right]\frac{mb}{(Gev)^4} \qquad (Ia')$$

From the analysis of the experimental data for $pp \to pX$ [4,5] one may obtain $\gamma_{PPR} \simeq 5$ and $\gamma_{RRP} \simeq 100$ with an accuracy up to 50%, i.e. at $t \to 0$ the Pomeron is connected with an accuracy up to 50% more weakly than the normal trajectories $\sigma^{\rho P}(M^2) \ll \sigma^{\rho R}(M^2)$. Hence it follows that at high but fixed S, and in the region $S^{3/2} \gg M^3 \gtrsim 5 \cdot 10^{-2} S$, the spectra will be described by the second term of (6) which essentially narrows the region of electromagnetic dominance. The change of the modes in (6) takes place at $M \simeq 5.5$.

In Fig.2, for illustration, it is plotted the region of the phase - where the dominance of the electromagnetic mechanism over the strong one must occur for the process $pp \to pX$ at S = 3600 Gev and $S \to \infty$. The sharp decrease of the dominance region for the case of real energies compared with the case of $S \to \infty$ is connected with the second term of (Ia) and with the presence of the factor $\left(1 - \frac{t_{min}}{t}\right)$ in the electromagnetic part. For the process $\pi p \to \pi X$, the influence of this factor is is weaker and $d\sigma_{st2}$ is about two times smaller, therefore, the electromagnetic effects become essential at much smaller values of S than in the case of $\bar{p}p \to pX$.

In the region under consideration,the contribution of interference takes the form

$$\frac{d\sigma_{int}}{dt\,dM^2} = \sum_{R=\rho,\omega} \frac{\alpha\,\Gamma_{\gamma R,p}(0)}{\sqrt{-t}\,\sqrt{SM^2}} \left(1-\frac{t_{min}}{t}\right)^{1/2} \qquad (4')$$

The contribution of (4') at $-tM^2 \leq (10^{-5} \div 10^{-6})\, S(1-\frac{t_{min}}{t})$ is comparable with the second term of (Ia').

If the experimental data at large M allow one to separate reliably the contribution of the background (RR,P) from that of the resonances (PP,R), then the corresponding electromagnetic terms ($\Gamma\Gamma$,P) and ($\gamma\Gamma$, R) appearing due to the separation of $\sigma_T^{\gamma p}(M^2) = \alpha\,\Gamma_{\gamma\gamma P}(0) + \frac{\alpha\,\Gamma_{\gamma\gamma R}(0)}{M^2}$ in (2a), will give Contribution to each of them, and the term ($\gamma\Gamma P$) will imitate the mechanism (PP,P) at low fixed -t. In the case of the separation of these contributions only by their asymptotic behaviour with respect to S, the above given asymptotic consideration is valid for the contribution which is independent of S at $S \to \infty$.

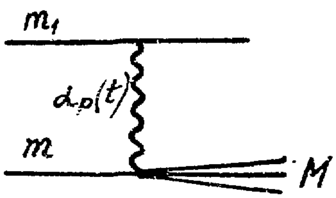
Let us note that the background can be studied in the process $pp \to \Delta X$ at large M, since in this case only ρ, A_2 and π exchanges are allowed and, therefore, we have the following formula:

$$\frac{d\sigma_{st2}}{dt\,dM^2} = \frac{1}{S^2} \sum_{R=\rho,A_2,\pi} \Gamma_{RR,p}(t)\left(\frac{S}{M^2}\right)^{2\alpha_R(t)} \underset{t\to 0}{\approx} \sum_{\rho,A_2} \frac{\Gamma_{RR,P}(0)}{S} + o\left(\frac{M^2}{S^2}\right) \qquad (6)$$

For the comparison we bring the corresponding electromagnet cross section:

$$\frac{d\sigma_{em}}{dt\,dM^2} = \frac{2\Gamma_{\Delta\to p\gamma}\,\sigma_T^{\gamma p}(M^2)}{\pi} \frac{M_\Delta^3}{(M_\Delta^2-m^2)^3} \frac{1}{M^2}\left[1+\left(1-\frac{t_{min}}{t}\right)^2\right] \qquad (7)$$

$\Gamma_{\Delta\to p\gamma} = 0.72$ Mev.



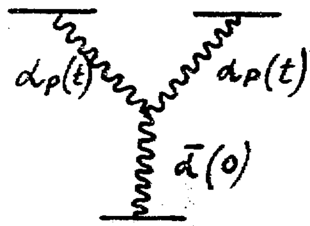
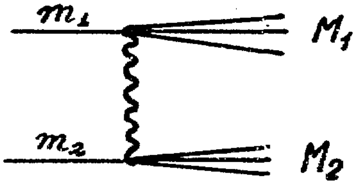


Fig. 1a



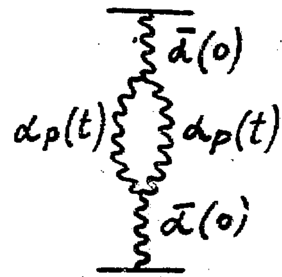
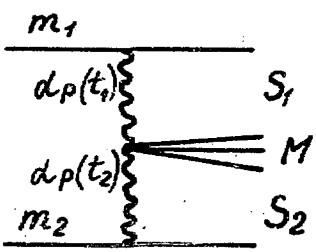


Fig. 1b



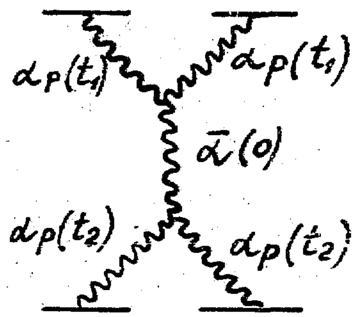


Fig. 1c.

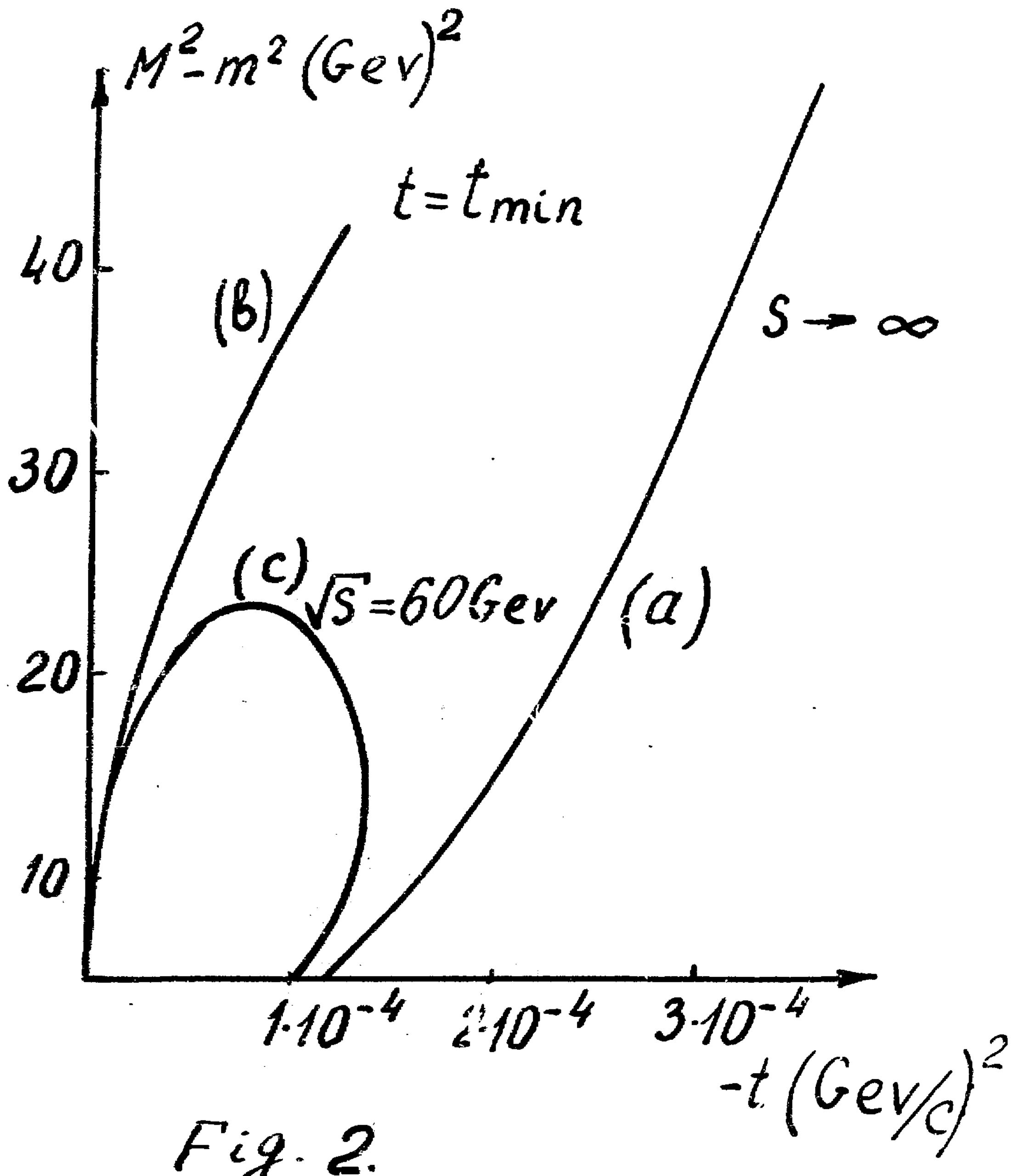


Fig. 2.

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