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QUARK-GLUON PLASMA AND HADRONIC PRODUCTION OF LEPTONS, PHOTONS AND PSIONS

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A B S T R A C T

Q C D calculations of the production rate in quark-gluon plasma and the account of space-time picture of hadronic collisions lead to estimates of dilepton mass M spectrum, $p_\perp$ distribution of $e^\pm$, $\mu^\pm$, γ, $\pi^\pm$ production cross section of charm, J/ψ, ψ' and χ mesons. For M, $p_\perp$ = 1 ÷ 5 GeV this theory describes data quite well, while for larger M, $p_\perp$ hard collisions dominate.

I. Introduction

Hadronic reactions, which take place at small and large distances, one treats on quite different theoretical ground. While the former are well described by parton model, based on asymptotic freedom of QCD, the latter are still discussed in more phenomenological way.

I should like to argue in this paper, that a very important intermediate region exists, namely reactions, taking place far from the collision point and not obeying the parton model, but at the same time treatable by perturbative QCD calculations. This is possible if momentum transfer involved (particle masses M, transverse momenta $p_\perp$ etc) are large in strong interaction scale $M, p_\perp \gtrsim 1$ GeV, but still are much smaller than the collision energy $\sqrt{s}$. Clear, this region increases with s and for ISR energies it only opens and contains $M, p_\perp = 1 \div 4$ GeV.

The most known example of such reaction is the production of dileptons, which deviates from Drell-Yan [1] model predictions by the factor of $10^{2\pm1}$ for $M = 1 \div 3$ GeV. Such extra dileptons were discussed by many authors, and the most clear qualitative explanation of their origin has been given by Bjorken and Weisberg [2] : Such pairs produced at later stages of collision when antiquarks are more numerous than in colliding hadrons; and they can interact repeatedly.

Much earlier, E.L.Feinberg [3] ascribed such pairs (and photons) to charge-current fluctuations in the framework of hydrodynamical model [4], and also stressed the importance of space-time aspect of the problem. But, because of the lack of any definite "hadronic matter" theory, he was not able to give the quantitative estimates of the cross sections.

Now, in view of recent theoretical development of QCD, it is highly probable that such matter is not a gas of hadrons but rather quark-gluon plasma, the properties of which are discussed in [5] (see also [6] on the nature of gas-plasma transition). As is shown below, the Q C D estimates lead to the conclusion, that at latest stages of the hadronic colli - sions the local [7] thermodynamical equilibrium is probably reached, but not at early stage, where the particles "go through" without scattering.

To give the idea let us say now, that temperatures we speak about for ISR do not exceed 0.5 GeV. This is rather low value and therefore production of heavy objects with $M,(p_\perp)$ of the order of 1 ÷ 5 GeV is suppressed by Boltsmann factor $\exp(-M/T)$. Such factor is absent for hard colli - sions and therefore one may think the low temperature plasma stage to be ineffective for production of such particles. Nevertheless, small Boltsmann factor is canseled by the large preexponent proportional to the space-time volume of the plasma stage, and so the discussed mechanism remains dominant up to $M, p_\perp$ = 4 ÷ 5 GeV. This conclusion is our main result and it is in agreement with more qualitative arguments of ref. [2,3].

The paper is structured as follows. The space-time picture of the collisions, corresponding to Feynman scaling [8] and Landau hydrodynamics [4] is discussed in chapter 2. The chapter 3 is devoted to dileptons due to reaction $q\bar{q} \to \mu^+\mu^-$. Chapters 4,5,6 are devoted to large $p_\perp$ leptons, photons and hadrons; in chapter 7 the total cross section for charm production is calculated; and that for charmonium states such as $J/\psi, \psi', \chi$ mesons is presented in chapter 8. Conclusions and final remarks are given in chapter 9.

2. The space-time picture of hadronic collisions

At Fig. 1 the main stages of the hadronic collision are shown. They are: (1) 2 structure function formation; (2) - hard collisions of constituents; (3) - final state interaction and (4) the propagation of free secondaries.

The phenomenological account for the first one and more explicite treatment for the second one is given by the parton

model. Our aim is to describe the stage (3), for which the plasma model is proposed. Therefore, these two approaches to particle production correspond to independend mechanisms, ta - king place in different space-time regions, so the predicted cross sections must be simply added.

It is better to begin our discussion of space-time picture from the stage (4), at which it is most simple. Indeed, secondaries are free and all come from the collision point $x\approx 0$, $t\approx 0$, so their velocities are connected with their coor - dinates x, t, by

$$V(x,t) \approx x/t \quad (1)$$

and their trajectories at Fig. 1 are straight lines. Of course, at stage (3) it is not so because of the particle interaction. But if the rapidity y distribution of secondaries is flat: dN/dy = const (the famous Feynman scaling), (1) remains valid: for any group of secondaries their interaction with left and right neighbours is identical and (in average) does not change their velocity. Such expansion without pressure gradient is analogous to that of the Universe, e.g. (1) reminds the Habble's low.

One can easily show (see e.g. (9)) that such scale inva - riant expansion sati sfies the relativistic hydrodynamical equations. Of course, this is not so for boundary conditions (plateau must have ends) and initial ones (expansion does not begin from a point) and therefore in more realistic approach deviations from scaling are inevitable. An example of it is the Landau theory [4]. Without going into details now we only comment that althogh initial conditions used in it are not satisfactory from various viewpoints [10], experimental data strongly support scaling violation of the type predic - ted by this theory [11,12]. Below we discuss Feynman scaling and Landau theory as too extreems. The question which one is more realistic is briefly discussed at the end of this chapter.

We asume the local thermal equilibrium of plasma at stage (3), which means that all its properties (including any reac - tion rate) are expressed through the single parameter, temperature T, being some function of x,t. This assumption, although approximate, *essentially* simplifies the problem.

In this paper we do not discuss the rapidity distribution of secondaries, but only their total production cross sections. So we integrate the production rate W_a of any particle "a" (per unite volume in unit time) over the whole space-time volume, occupied by the plasma stage:

$$\sigma_a = \sigma_{in} \int_{(3)} d^4x \, W_a = \sigma_{in} \int_{T_i}^{T_f} dT \, W_a(T) \Phi(T) \quad (2)$$

where σ_{in} is the inelastic cross section; and $\Phi(T)$ is the temperature distribution function defined as

$$\Phi(T) = \int_{(3)} d^4x \, \delta(T(x,t) - T) \quad (3)$$

which must be the same for any reaction. Moreover, it turns out that this function is also the same for both theories considered, namely:

$$\Phi(T) \simeq A \cdot T^{-7} \quad (4)$$

For scaling expansion this result can easily be found by transfer to comoving frame: $\tau = (t^2 - x^2)^{1/2}, \alpha = \frac{1}{2} \ln\left(\frac{t-x}{t+x}\right)$ in which the particle density is $n \sim T^3 \sim 1/\tau$. Landau solution [4] also leads to (4) (up to log terms), and it was also checked numerically that (4) survives in more realistic case $c^2 = 0.2$ and three dimensional expansion (see e.g. [12,13]), which will be published elswere. The only difference of these two approaches is the energy dependence of normalization factor A in (4): $A \sim \ln S$ and $S^{1/2}$ respectively.

Now let us discuss the interaction limits in (2), the initial T_i and final T_f temperatures. The latter one, is $T_f \sim m_\pi$, as known from spectra of secondary hadrons.

The T_i depends on the relaxation rate in hadronic matter and in fact is the only parameter of the discussed theory.

In Feynman parton approach local equilibrium means that local longitudinal momenta spread becomes of the order of that of transverse ones, so T_i is connected with $\langle p_\perp \rangle$ of partons: $T_i \simeq \langle p_\perp \rangle / 2$. As is seen from data, $\langle p_\perp \rangle = 1 \div 1.4$ GeV and $T_i = 0.5 \div 0.7$ GeV independent on collision energy.

In Landau theory the initial temperature is determined from the normalization to energy density at the collision moment and is estimated as

$$T_i = \left(\frac{45 \, S}{4\pi^3 \tau_0^3 m \varkappa} \right)^{1/4} \simeq 0.12 \left(\frac{S}{4m^2} \right)^{1/4} (GeV) \quad (5)$$

where we have substituted the typical values of hadron dimensions $\tau_0 \sim 1/m_\pi$ and the effective number of degrees of freedom $\varkappa = 30 \cdot \varepsilon / \pi^2 T^4$ (ε is energy density). For ideal gas of colored quarks, antiquarks and gluons $\varkappa \simeq 40$, but with the account of interaction one finds for $T \sim 0.5$ GeV $\varkappa \sim 20$ [5]. Note, that because of large $\varkappa$ T_i is essentially smaller than original Landau estimates. From (5) we see that for ISR energies $T_i \sim 0.5$ GeV which is close to that in parton model and therefore one has to wait for higher energies in order to separate them.

The value of the normalization factor A in (4) can also be estimated from initial volume value in Landau theory with the result:

$$A = 4\pi \tau_0^4 m^2 T_i^6 / S \simeq 0.02 (S/4m^2)^{1/2} (GeV^2) \quad (6)$$

Apart from reactions discussed below, this estimates can be checked by the total hadronic multiplicity and is consistent with data.

For definitness, in comparison with data to be given below at Fig. 2-4 we use T_i, A values (5), (6), although in fact they are neither accurate enough nor taken at high enough energy for the real check of such energy dependence.

After presenting these model considerations, we proceed to brief discussion of more direct way to estimate T_i, which is an explicite solution of partonic kinetics during the collision. We know the initial conditions, the proton structure functions, but we don't know properly the scattering cross sections for small momenta transfer, for which $\alpha_s \sim 1$. We only do know, that large momentum transfer reactions have small cross sections and therefore one cannot expect T_i to be large. So, at large enough energy partons really "go through", as proposed by Feynman, rather than "glued together", as proposed by Fermi and Landau. Such arguments are, of course, widely known. Nevertheless the point is, that this may happen at high energies $\sqrt{s} \gtrsim 10^2$ GeV because of two reasons. The rather first one is more evident: gluons and see quarks carry only small part of the total momentum. The second one is less trivial: the large number of final states (due to color) lead to rather large numerical factors in Q C D cross sections. To give an example, the large angle gluon-gluon scattering $d\sigma/dt \approx 100\,\alpha_s^2/s^2$ [14] and due to this the mean free path of gluon in plasma is $\ell_0 \sim (30\alpha_s^2 T)^{-1}$. Stricly speaking, all this is valid for $\alpha_s \ll 1$, but nonperturbative effects (instantons etc.) probably only increase the relaxation rate. So one is lead to the conclusion that up to ISR energies relaxation length ℓ_0 may be rather small and the expansion picture is more close to the Landau version, explaning its agreement with data.

To prove or reject these speculations one can with the help of reactions which we discuss in next chapters.

3. Dileptons

We begin with the calculation of the production rate $W_{\mu^+\mu^-}$ - in unite volume of plasma of u, d, s quarks and antiquarks due to reaction $q\bar{q} \to \mu^+\mu^-$ (e^+e^-), which is equal to

$$W_{\mu^+\mu^-} = \int |\langle M\rangle|^2 (2\pi)^4 \delta(\Sigma p_i)\, n(\varepsilon_1)\, n(\varepsilon_2) \prod_{i=1}^{4} \frac{d^3p_i}{\varepsilon_i (2\pi)^3} \qquad (7)$$

where $\langle M\rangle$ is the standard matrix element; p_1, p_2, p_+, p_- are momenta of $q, \bar{q}, \mu^+, \mu^-$ respectevely; $n(\varepsilon) = [\exp(+\varepsilon/T)+1]^{-1}$, is the Fermi occupation factor. As it was found earlier [5], for massless quarks and leptons

$$W_{\mu^+\mu^-} = \frac{\pi\alpha^2 T^4}{108} \qquad (8)$$

but now we are interested in the spectra of dilepton masses $M = (p_+ + p_-)^2$. In the case $M \gg T$ one can neglect interference effects so, that $n(\varepsilon_1)\, n(\varepsilon_2) \approx \exp(-\frac{\varepsilon_1+\varepsilon_2}{T})$ and make use of energy conservation: $\varepsilon_1+\varepsilon_2 = \varepsilon_+ + \varepsilon_-$. Then integration can be done, leading to the following expression in terms of the cross section of the <u>inverse</u> process

$$W_{\mu^+\mu^-} = \int \frac{d^3p_+}{\varepsilon_+ (2\pi)^3} \frac{d^3p_-}{\varepsilon_- (2\pi)^3} \exp\left(-\frac{\varepsilon_+ + \varepsilon_-}{T}\right) 2M^2 \sigma(\mu^+\mu^- \to q\bar{q}) \qquad (9)$$

This expression is the manifistation of the detailed balance: the rate of direct and inverse reactions in equilibrium are equal. In reality the lepton's density is small and inverse precess is negligible, so (9) is just the convenient form for the probability of the direct one, to be systematically used below. From (9) one get then

$$\frac{dW_{\mu^+\mu^-}}{dM} = \frac{\sqrt{2}\,\alpha^2 (MT)^{3/2}}{3\pi^{5/2}} \exp\left(-\frac{M}{T}\right) \qquad (10)$$

and intergating over space-time (2-4), we have the cross section of the dilepton production with a given mass:

$$\frac{d\sigma}{dM} = \frac{\sqrt{2}\,\alpha^2 \sigma_{in} A}{3\pi^{5/2} M^3} \left[\Gamma\left(\frac{9}{2}, \frac{M}{T_i}\right) - \Gamma\left(\frac{9}{2}, \frac{M}{T_f}\right) \right] \qquad (11)$$

Here function $\Gamma(\alpha, x)$, is the incomplete gamma function.

At Fig.2 the comparison is made of (11) with the data [15, 16] at $\sqrt{s}$ = 22, 28 GeV and the parton model [1]. One can see that for $M \lesssim 4$ GeV data agrees with (11), while for larger M the Drell-Yan mechanism works well. Let us add that at $M \sim 4$ GeV the change is observed in various features of the phenomenon: atomic number dependence; the value of $\langle p_\perp\rangle$ etc, indicating the change of the production mechanism. Note also, that higher masses are connected with higher temperatures, then $\langle p_\perp^2\rangle$ must grow with M as $2MT$. Data support such M - dependence, but the coefficient for initial period is somehow smaller than T. This observation means the incomplete mixing for $T \sim T_i$ of the transverse and longitudinal momenta of quarks in plasma. Similar behaviour is demonstrated by $\langle p_\perp^2\rangle$ of psions. Probably the best way to study this question is the Θ^* distribution of dileptons (Θ^* is the C.M. angle between lepton momenta and that of the initial hadron), but now such data are not good enough.

4. Direct leptons

In this chapter we discuss experiments with the detection of the so called direct leptons, not being the product of π, K decay. As later studies have shown, they can not be explained completely by the decay of $\rho, \omega, \varphi, \Upsilon$ mesons [17].

The mechanism of direct lepton production we are going to discuss is the same as that for dileptons, namely $q\bar{q} \to \mu^+\mu^-$. Integrating (9) over all variables except the transverse momentum one has

$$\frac{dW}{dp_\perp^2} = \frac{\alpha^2 T^2}{18\pi} K_0\left(\frac{p_\perp}{T}\right) \qquad (12)$$

and therefore the cross section for production of any type of leptons $\mu^+ = \mu^- = e^+ = e^-$ is

$$\frac{d\sigma}{dp_\perp^2} = \frac{\alpha^2 \sigma_{in} A}{18\pi p_\perp^4}\left[F\left(\frac{p_\perp}{T_i}\right) - F\left(\frac{p_\perp}{T_f}\right)\right] \qquad (13)$$

$$F(x) = \int_x^\infty dz \cdot z^3 K_0(z)$$

The corresponding curves for $\sqrt{s}$ = 24, 52 GeV are compared with data [18, 19] at Fig.3. The dashed curve below is the parton model prediction, it lies much below the data. As for (13), it agrees quite well with FNAL data, but clearly overestimates the cross section for largest $p_\perp$ at ISR. The reason is again the incomplete mixing at the initial stage.

5. Direct photons

The so called direct photons are those which are not due to π^0, η decays, they were first observed quite recently [20]. In the lowest orders in QCD the corresponding reactions are

$gq \to \gamma q$; $g\bar{q} \to \gamma\bar{q}$; $q\bar{q} \to \gamma g$, discussed in the framework of parton model in [21] . The γ production rate by the plasma can be written as

$$W(gq \to \gamma q) = \int \frac{d^3p_\gamma \, d^3p_q \cdot 18}{(2\pi)^6 \, \varepsilon_\gamma \, \varepsilon_q} \exp\left(-\frac{\varepsilon_\gamma + \varepsilon_q}{T}\right) s\,\sigma(s) \qquad (14)$$

where $\sigma(s)$ is the cross section of the <u>inverse</u> process, as in (9). For other channels expressions are similar and using

$$\sigma(\gamma q \to gq) = \frac{16\pi\alpha\alpha_s}{81 s} \ln\left(\frac{s}{4m^2}\right) ; \quad \sigma(\gamma g \to q\bar{q}) = \frac{4\pi\alpha\alpha_s}{3 \cdot s} \ln\left(\frac{s}{4m^2}\right) \qquad (15)$$

one has

$$\frac{dW_\gamma}{dp_\perp^2} = \frac{10\alpha\alpha_s}{27\pi} \ln\left(\frac{s}{4m^2}\right) T^2 K_0\left(\frac{p_\perp}{T}\right) \qquad (16)$$

so the ratio $\gamma/\mu^+ = 20\alpha_s \ln(p_\perp T/m_0^2)/3\alpha \approx 600^*$ and

$$\frac{d\sigma_\gamma}{dp_\perp^2} = \frac{10\alpha\alpha_s \sigma_{in} A}{27\pi p_\perp^4}\left[F\left(\frac{p_\perp}{T_i}\right) - F\left(\frac{p_\perp}{T_f}\right)\right] \qquad (17)$$

The corresponding curve marked " γ " at Fig.3 is in rather nice agreement with data [20] , but this point needs further experimental studies.

6. Evaporating hadrons

In contrast to leptons and photons, hadrons are "dosed up" by their interaction and normally are confined in the system up to the final breaking to secondaries at low $T_f \sim m_\pi$ [4]. This idea is very well supported by low $p_\perp \leq$ 1 GeV data, see e.g. [7a] . For higher $p_\perp$ the evaporation mechanism at earlier stages with higher temperatures was proposed [7] . In this chapter we discuss two possible approaches to such process:

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* Stricly speaking $\alpha_s(s) \ln s \sim 1$ and further terms in this parameters must be included. So (17) is only an order of magnitude estimate.

1) the evaporation of pions; 2) the leakage of quarks or gluons with large $p_\perp$, then decaying into hadronic jet. Both give the particle flux at the surface

$$dj_a = \frac{g_a T}{4\pi^2} \exp\left(-\frac{p_\perp}{T}\right) dp_\perp^2 \qquad (18)$$

where $g_a = 1$ in the first case, while in the second one has to know the jet structure function. If one makes the usual guess $G(q\to\pi) \sim (1-x)$; $G(g\to\pi)\sim(1-x)^3$ he finds at fixed $p_\perp$ π^+/(quark jet) $\sim 2\cdot 10^{-2}$; π^+/(gluon jet) $\sim 10^{-3}$ and with the account of quark colors and flavors it finally gives in (18) $g_a \sim 0.2$. From (18), neglecting transverse expansion and taking the system to be the cylinder of radius r_0 $(\sim 1/m_\pi)$, one finds

$$\frac{d\sigma}{dp_\perp^2} = \frac{g_a \sigma_{in} A}{16\pi^2 r_0 p_\perp^3}\left[F\left(\frac{p_\perp}{T_i}\right) - F\left(\frac{p_\perp}{T_f}\right)\right] \qquad (19)$$

$$F(x) = \exp(-x)\left[x^4 + 4x^3 + 12x^2 + 24x + 24\right]$$

the corresponding curves are shown at Fig. 3 together with data for $\pi^+ = \pi^-$ [22] and $= \pi^0$ [20]. The dashed curve represents the hard scattering QCD calculations [14a]. Comparing data with theoretical predictions, one may conclude that these theories work below and above $p_\perp \sim 4$ GeV, respectevely.

It is important to note, that at $p_\perp \sim 4$ GeV many features of the phenomenon are respectevely changed: 1) π^+/π^- ratio jumps from 1 to 2 (in pp collisions); $\bar{p}/\pi$ ratio stop to grow (in agreement with temperature growth [7a] and falls rapidly; 3) atomic number dependence A^α give $\alpha > 1$ for $p_\perp \sim$ ~ 4 GeV, but $\alpha = 1$ at higher $p_\perp$ etc. All these facts support the idea about the change of the production mechanism at $p_\perp \sim 4$ GeV.

7. Charmed quarks

The plasma of light quarks (u, d, s) created in the collision, produces with some rate pairs of heavy quarks $c\bar{c}$. Since their total mass is ~ 3 GeV and in view of our results presented above, one may think that such pairs are produced mainly at plasma stage. This is not so for new quarks which probably compose the $\Upsilon(9.5)$ particle and we are not going to discuss its production.

Using again the cross section of the inverse process we can write

$$W_{c\bar{c}} = \int \frac{6\, d^3p_1\, n(\varepsilon_1)}{(2\pi)^3} \frac{6\, d^3p_2\, n(\varepsilon_2)}{(2\pi)^3} v\sigma(c\bar{c}\to gg) \qquad (20)$$

and analogous expression for $c\bar{c} \to q\bar{q}$.

Since in typical conditions $2m_c \gg T$, c quarks are nonrelativistic and we need to know only the small relative velocity limit $v\to 0$ of the product $v\sigma$, which is

$$v\sigma(c\bar{c}\to gg) = \frac{49\pi\alpha_s^2}{81 m_c^2}; \quad v\sigma(c\bar{c}\to q\bar{q}) = \frac{2\pi\alpha_s^2}{3m_c^2} \qquad (21)$$

Note that in [23,24] three-gluon diagram was not included. The value of α_s in (21) is $\alpha_s(m_c^2) = 0.2$ (see e.g. [25]). Finally, the total cross section for $c\bar{c}$ production is

$$\sigma_{c\bar{c}} = \frac{103\,\alpha_s^2\,\sigma_{in} A}{72\pi^2 m_c^2}\left(1 + \frac{2m_c}{T_i} + \frac{2m_c^2}{T_i^2}\right)\exp\left(-\frac{2m_c}{T_i}\right) \qquad (22)$$

shown at Fig. 4 together with the parton model contribution [23,24,26]. One may see that at FNAL energies the predictions are very close in magnitude, but in fact the resulting final states must be different: mainly $D\bar{D}$ pairs in parton mechanism and bound states (psions)in plasma one. Experimentally only some / upper bounds for such cross sections are known, which do not contradict yet to these estimates.

8. Psions

In this chapter we discuss the production of charmonium bound states, or psions.

The general expression for the production rate may be integrated over momenta of initial particles and put into universal form:

$$W_a = \tilde{\Gamma}_a (2S_a + 1) \left(\frac{m_a T}{2\pi} \right)^{3/2} exp\left(-\frac{m_a}{T}\right) \quad (23)$$

where we have used $m_a \gg T$; = 25a + 1, S_a is the spin value; $\tilde{\Gamma}_a$ is the decay width of the particle <u>in plasma.</u> This quantity differs from that in vacuum (apart from interaction corrections) just by interference of the decay products (quarks and gluons) with those of plasma. In the case $T \ll m_a$ this effect is small, but for J/ψ mesons, decaying into three gluons, for T = 0.4 ÷ 0.6 GeV the ratio $\tilde{\Gamma}/\Gamma$ = 1.2 ÷ 1.8. For χ mesons such corrections (induced decay) are not larger than 10%.

Since c quark density is small, recombination $c\bar{c} \to g\psi$ is negligible. It means that the inverse process $g\psi \to c\bar{c}$ ("photoeffect") must not be included in $\tilde{\Gamma}$ in (23). Never - theless, this process is interesting in order to know what part of psions, created in plasma, survive during the expansion. The simple nonrelativistic estimate $\sigma(g\psi \to c\bar{c}) \sim \pi \alpha_s a^2$, $a \sim 1/(\alpha_s m_c)$ leads to $\sigma \sim$ 3 mb and therefore this pro - cess is negligible. Interesting, that experimental absorbtion at nuclei corresponds to σ = 3.5 ± 1.0 [27].

As is well known, J/ψ and ψ' mesons decays into three gluons, and therefore the inverse process is [3] collisions. Of course, it is quite possible in finite density plasma, but not in the framework of the parton model. The corresponding cross section of J/ψ and ψ' is

$$\sigma_\psi = \frac{3 \tilde{\Gamma}_\psi \, \sigma_{in} A}{(2\pi)^{3/2} \, m_\psi^3} \, \Gamma\left(\frac{9}{2}, \frac{m_\psi}{T_i}\right) \quad (24)$$

According to [23,26] for J/ψ the additional channel is the following reaction: $2g \to \chi \to \gamma + J/\psi$. Its cross section can also be calculated by (24) with the obvious change $3 \to (2S_\chi + 1)$ $\tilde{\Gamma}_\psi \to \tilde{\Gamma}(\chi \to \gamma + J/\psi)$. Although experimentally $\tilde{\Gamma}(\chi \to \gamma + J/\psi)$ are not known, but they can be estimated theoretically [25] : for χ mesons of the type $3P_0$ /3414 Mev/; $3P_1$, (3508); $3P_2$ (3552) they are equal to 0.1 ÷ 0.2; 0.3 ÷ 0.4; 0.3÷0.5 Mev.

At Fig. 4 the corresponding curves for J/ψ and ψ' mesons are compared with data [28]. For J/ψ the sum of two mechanisms are shown, their ratio $(3g \to J/\psi)/(\chi \to \gamma + J/\psi)$ = 0.2÷0.3. Recently [29] the accomponing γ has been observed with the conclusion that $\chi \to \gamma + J/\psi$ gives 43 ± 21% of the cross section. This contradicts strongly to parton model, where there is no competing process, while in the framework of our model it means only that $3g \to J/\psi$ process is a factor 2 underestimated, which is quite inside the general uncertainties.

9. Conclusions

The main conclusion is that the proposed theory is able to give very simple and consistent describtion of various reaction in $M, p_\perp$ = 1 ÷ 5 GeV region. It is also very important, that just above this region of $M, p_\perp$ the hard scattering was known to dominate. Changes in various characteristics at $M, p_\perp \sim$ 4 GeV also support the change of production mechanism in this region.

The second comment is that ingredients of this theory like T_i value, $\Phi(T)$ dependence etc. turns out to be very stable in sense that various theoretical approaches lead to the same result. This is not so for rapidity and energy dependence of spectra, which are therefore interesting for model separation, in particular for the check of the predicted transition from Landau space-time picture to Feynman one at $\sqrt{s} \sim 10^2$ GeV.

There are open questions concerning the degree of reaching local thermal equilibrium, which may be checked in many ways experimentally. From the theoretical point of view, further progress in QCD needed to treat the parton kinetics properly.

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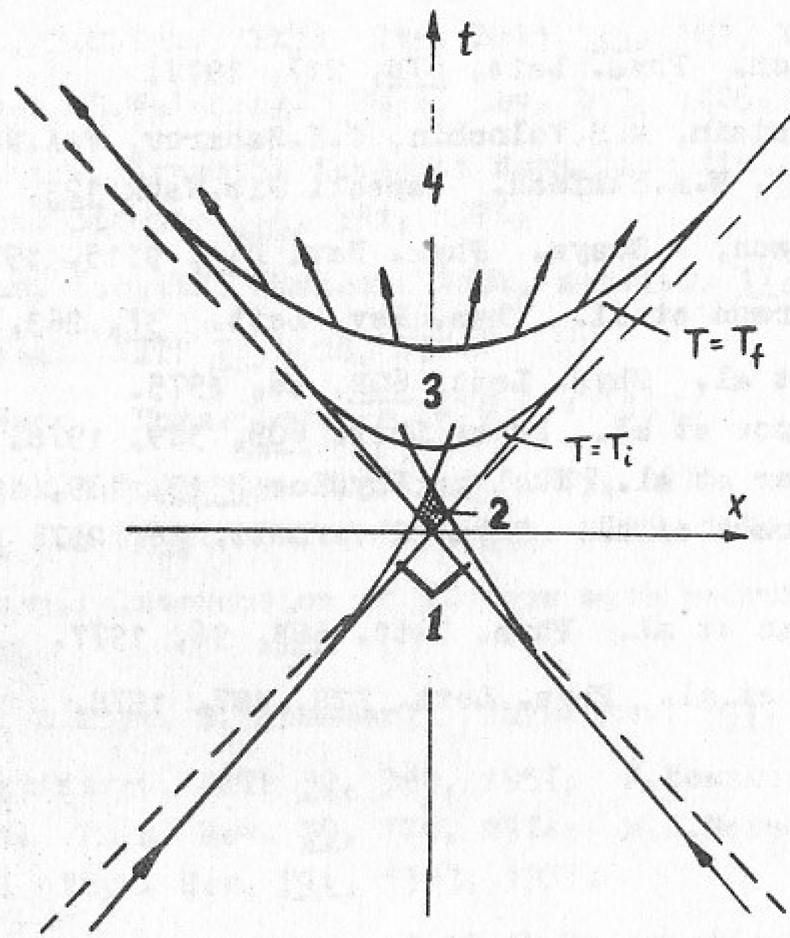

Fig. 1. The space-time picture of hadronic collisions, pro - ceeding via the following stages: 1) structure function formation; 2) hard collisions; 3) final state interaction; 4) free secondaries.

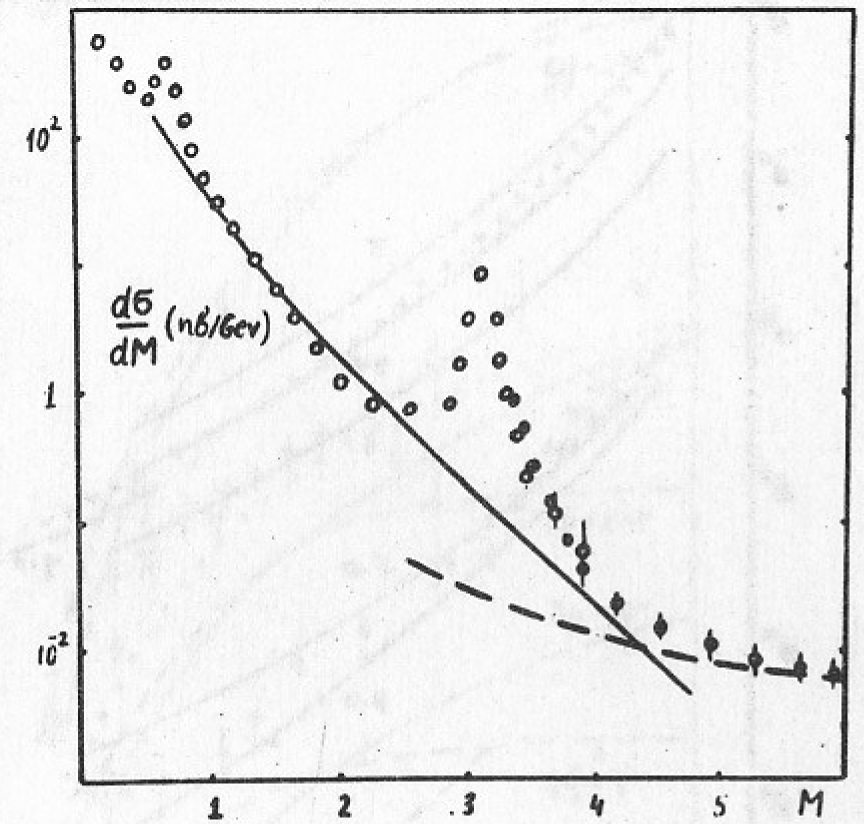

Fig. 2. Dilepton mass spectrum $d\sigma/dM$ (nb/GeV.Nucleon) versus M(GeV). Solid and dashed lines correspond to (11) and [1] (with color), open and black points to data [15 , 16].

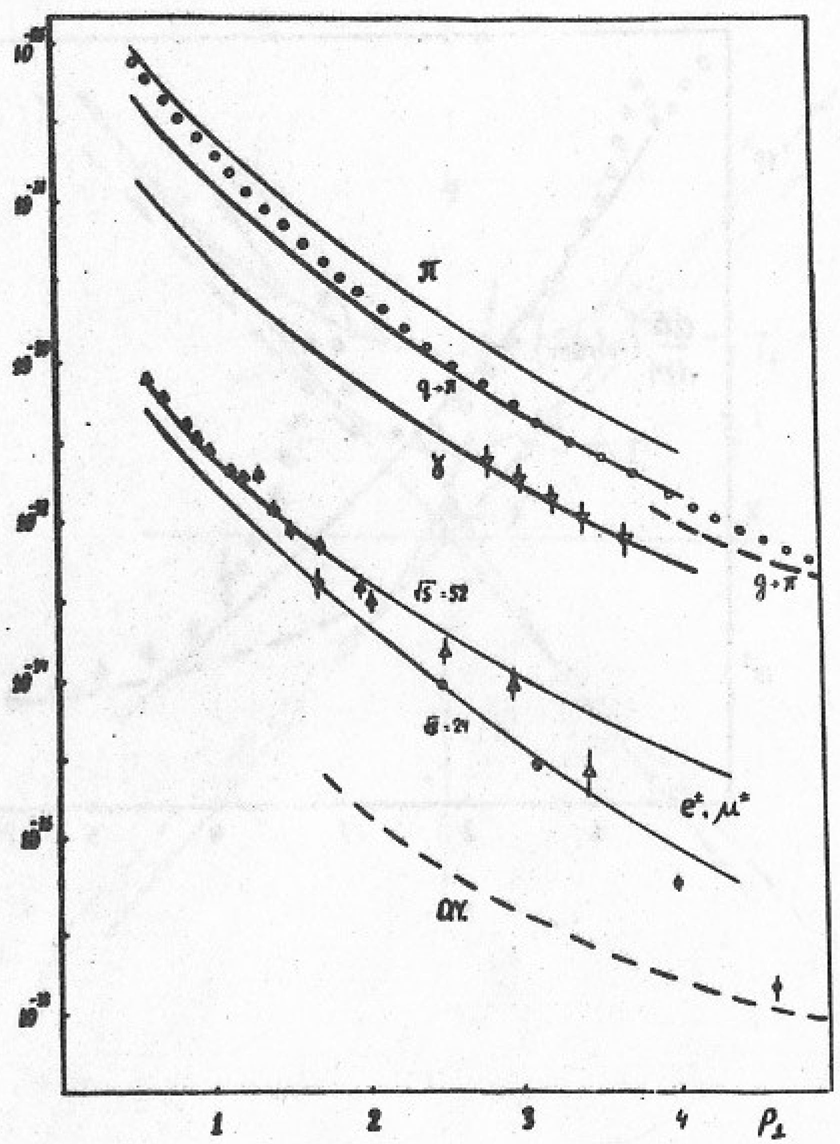

Fig. 3. $Ed^3\sigma/dp^3$ ($cm^2 \cdot c^3/GeV^2$) at $y = 0$ versus $p_\perp$ (GeV/c) for: o - $\pi^+ = \pi^-$ [22], $= \pi^0$ [20] at $\sqrt{S}$ = 53 GeV; ∇ - direct γ [20] at $\sqrt{S}$ = 53 GeV; Δ - direct e^- [19] at $\sqrt{S}$ = 53 GeV; ● - direct μ^+ at $\sqrt{S}$ = 24 GeV [18]. Solid lines are described in the text; the dashed one marked "D.Y." corresponds to [1] and that marked " $g \to \pi$ " " to [14 a], they both respenst hard scattering models.

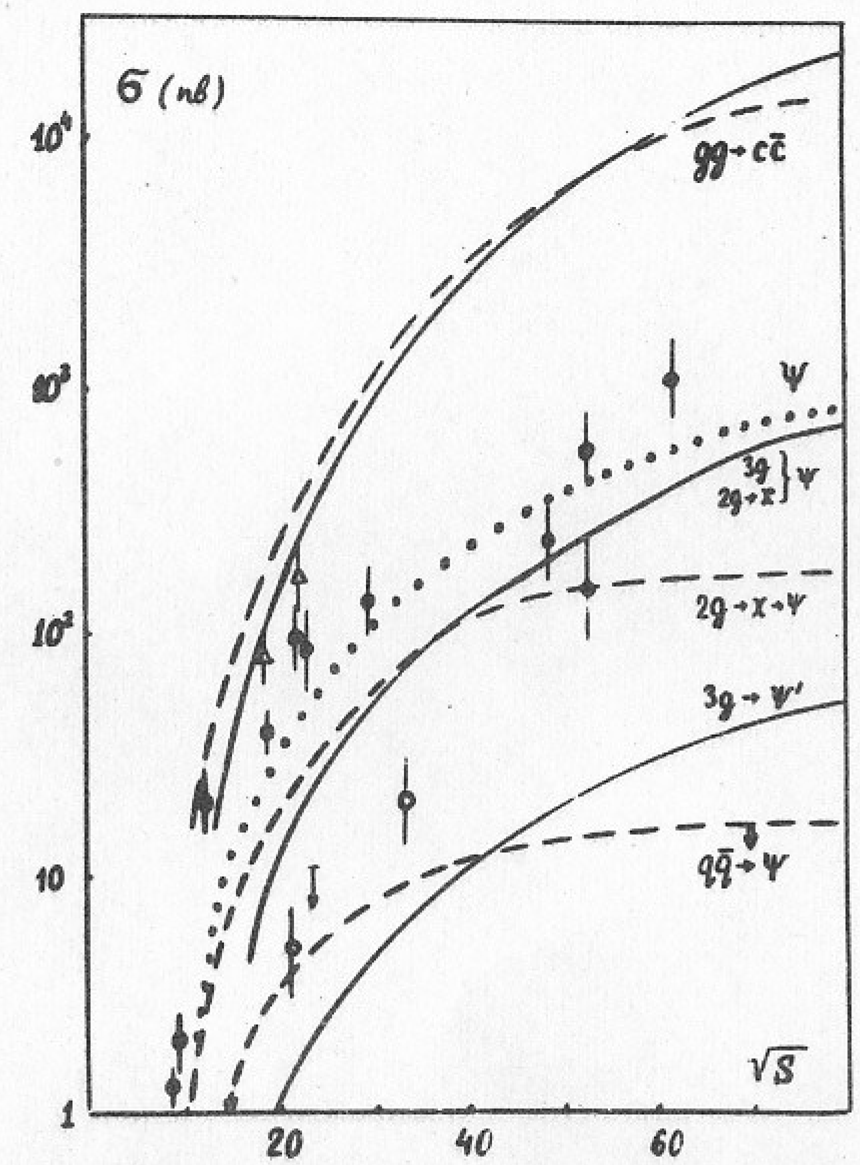

Fig. 4. The cross section σ (nb) of chaı and psion production versus $\sqrt{S}$ (GeV). Points [28] mean ; ● - J/ψ in pp ; Δ - J/ψ in πp ; o - ψ' in pp . Solid lines correspond to the next, the dashed ones-to parton model [26] . The dotted one is their sum for $pp \to J/\psi$. The curve marked $q\bar{q} \to \psi$ is the upper bound in parton model based on the restrictions from the total width of J/ψ .

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