On the hadron production from the quark–gluon plasma phase in ultra–relativistic heavy–ion collisions

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Abstract

We describe the quark gluon plasma (QGP) as a thermalized quark–gluon system, the thermalized QGP phase of QCD. The hadronization of the thermalized QGP phase is given in a way resembling a coalescence model with correlated quarks and anti–quarks. The input parameters of the approach are the spatial volumes of the hadronization. We introduce three dimensionless parameters $C_M$, $C_B$ and $C_\bar{B}$ related to the spatial volumes of the production of low–lying mesons ($M$), baryons ($B$) and antibaryons ($\bar{B}$). We show that at the temperature $T = 175$ MeV our predictions for the ratios of multiplicities agree good with the presently available set of hadron ratios measured for various experiments given by NA44, NA49, NA50 and WA97 Collaborations on Pb+Pb collisions at 158 GeV/nucleon, NA35 Collaboration on S+S collisions and NA38 Collaboration on O+U and S+U collisions at 200 GeV/nucleon.

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1 Introduction

Recently [1] we have suggested some kind of a coalescence model [2] for the description of the hadronization from the quark–gluon plasma (QGP) phase of QCD considered as a thermalized quark–gluon system [3] at high densities and temperature in which quarks, antiquarks and gluons being at the deconfined phase collide frequently each other. There is a belief [4] that the QGP phase of the quark–gluon system can be realized in ultra–relativistic heavy–ion collision ($E_{\text{cms}}/\text{nucleon} \gg 1 \text{ GeV}$) experiments.

At very high energies of heavy–ion collisions the quark–gluon system is composed from highly relativistic and very dense quarks, antiquarks and gluons. By virtue of the asymptotic freedom the particles are almost at liberty and due to high density collide themselves frequently that leads to an equilibrium state. If to consider such a state as a thermalized QGP phase of QCD, the probabilities of light massless quarks $n_q(\vec{p})$ and light massless antiquarks $n_{\bar{q}}(\vec{p})$, where $q = u, d$, to have a momentum $p$ at a temperature $T$, can be described by the Fermi–Dirac distribution functions [3,5]:

$$n_q(\vec{p}) = \frac{1}{e^{-\nu(T)} + p/T + 1}, \quad n_{\bar{q}}(\vec{p}) = \frac{1}{e^{\nu(T)} + p/T + 1},$$

where a temperature $T$ is measured in MeV, $\nu(T) = \mu(T)/T$, $\mu(T)$ is a chemical potential of the light massless quarks $q = u, d$, depending on a temperature $T$ [5]. A chemical potential of light antiquarks amounts to $-\mu(T)$. A positively defined $\mu(T)$ provides an abundance of light quarks with respect to light antiquarks for a thermalized state [1,4]. A chemical potential $\mu(T)$ is a phenomenological parameter of the approach which we would fix below [1].

The probability for gluons to have a momentum $\vec{p}$ at a temperature $T$ is given by the Bose–Einstein distribution function

$$n_g(\vec{p}) = \frac{1}{e^{p/T} - 1}.$$ (1.2)

Since a strangeness of the colliding heavy–ions amounts to zero, the densities of strange quarks and antiquarks should be equal. The former implies a zero–value of a chemical potential $\mu_s = \mu_{\bar{s}} = 0$. In this case the probabilities of strange quarks and antiquarks can be given by

$$n_s(\vec{p}) = n_{\bar{s}}(\vec{p}) = \frac{1}{e^{\sqrt{\vec{p}^2 + m_s^2}/T} + 1},$$ (1.3)

where $m_s = 135 \text{ MeV}$ [6] is the mass of the strange quark and antiquark. The value of the current $s$–quark mass $m_s = 135 \text{ MeV}$ has been successfully applied to the calculation of chiral corrections to the amplitudes of low–energy interactions, form factors and mass spectra of low–lying hadrons [7] and charmed heavy–light mesons [8]. Unlike the massless antiquarks $\bar{u}$ and $\bar{d}$ for which the suppression is caused by a chemical potential $\mu(T)$, the strange quarks and antiquarks are suppressed by virtue of the non–zero mass $m_s$.

In Ref. [1] we have supposed that a chemical potential $\mu(T)$, a phenomenological parameter of the description of the QGP state as a thermalized quark–gluon system at a temperature $T$, is an intrinsic characteristic of a thermalized quark–gluon system. Thereby,
if the QGP is an excited state of the QCD vacuum, so a chemical potential should exist not only for ultra–relativistic heavy–ion collisions. Quark distribution functions of a thermalized quark–gluon system at a temperature $T$ should be characterized by a chemical potential $\mu(T)$ for any external state and any external conditions. Since any state of a thermalized system is closely related to external conditions, in order to obtain $\mu(T)$ we need only to specify the external conditions of a thermalized quark–gluon system the convenient for the determination of $\mu(T)$.

Indeed, it is well–known [5] that the Helmholtz free energy $F(T, V, N)$ defining the partition function $Z(T, V, N)$, $F(T, V, N) = - T \ln Z(T, V, N)$ which plays a central role in studying thermalized systems, is nothing more than a work for an isothermic process. Therefore, by producing external conditions keeping $T = \text{const}$ and measuring a work one can get a full information about the Helmholtz free energy $F_{\text{exp}}(T, V, N)$. Then, in terms of this Helmholtz free energy $F_{\text{exp}}(T, V, N)$ one can obtain the partition function $Z_{\text{exp}}(T, V, N)$ which can be applied to the description of the thermalized system at any $T$.

Following this idea we have fixed the chemical potential $\mu(T)$ in the form [1]:

$$\frac{\mu(T)}{\mu_0} = \left[ \frac{1}{2} + \frac{1}{2} \sqrt{1 + \frac{4\pi^6}{27} \left( \frac{T}{\mu_0} \right)^6} \right]^{1/3} - \left[ -\frac{1}{2} + \frac{1}{2} \sqrt{1 + \frac{4\pi^6}{27} \left( \frac{T}{\mu_0} \right)^6} \right]^{1/3}. \quad (1.4)$$

In the low–temperature limit $T \to 0$ we get

$$\mu(T) = \mu_0 \left[ 1 - \frac{\pi^2 T^2}{3 \, \mu_0^2} + O(T^6) \right]. \quad (1.5)$$

where $\mu_0 = \mu(0) = 250 \, \text{MeV}$ is a chemical potential at zero temperature [1]. The $T$–dependence of a chemical potential given by Eq.(1.3) differs by a factor $1/4$ from the low–temperature behaviour of a chemical potential of a thermalized electron gas [9]. The former is caused by the contribution of antiquarks.

In the high–temperature limit $T \to \infty$ a chemical potential $\mu(T)$ defined by Eq.(1.4) drops like $T^{-2}$:

$$\mu(T) = \frac{\mu_0^3}{\pi^2} \frac{1}{T^2} + O(T^{-7}). \quad (1.6)$$

A chemical potential drops very swiftly when a temperature increases. Indeed, at $T = 160 \, \text{MeV}$ we obtain $\mu(T) \simeq \mu_0/4$, while at $T = \mu_0$ a value of a chemical potential makes up about tenth part of $\mu_0$, i.e. $\mu(T) \simeq \mu_0/10$. This implies that at very high temperatures the function $\nu(T) = \mu(T)/T$ becomes small and the contribution of a chemical potential of light quarks and antiquarks can be taken into account perturbatively. This assumes in particular that at temperatures $T \geq \mu_0 = 250 \, \text{MeV}$ the number of light antiquarks will not be suppressed by a chemical potential relative to the number of light quarks.

In our approach the multiplicities of hadron production we define in terms of quark and anti–quark distributions functions in a way similar to a simple coalescence model [2] but for correlated quarks and anti–quarks. Indeed, in a coalescence model quarks and anti–quarks are uncorrelated [2]. This allows to introduce separately the number of light quarks $q$ and light anti–quarks $\bar{q}$ and the number of strange quarks $s$ and strange anti–quarks $\bar{s}$.
The factor 3 corresponds to the number of quark color degrees of freedom, $\frac{1}{3}N_f$ coupling constants of the $q\bar q$ of hadrons, $F$ we define as the relative momenta of order $T$ are coalesced into a meson with a 3-momentum $\vec{q}$ at a temperature $T$. Since the main contribution to the integrals comes from the relative momenta of order $p \sim T$, so that quark and anti–quarks coalesce at relative momenta of order $p \sim T$. This agrees with the order of transversal momenta of hadrons, $q_\perp \sim 2 \div 3 T$, coupled in the center of mass frame of heavy–ion collisions. The factor 3 corresponds to the number of quark color degrees of freedom, $M_K = 500$ MeV, $F_K = 160$ MeV, $M_\pi = 140$ MeV and $F_\pi = 131$ MeV are the masses and the leptonic coupling constants of the $K$ and $\pi$ mesons, respectively [10]. The dimensionless parameter $C_M$ is a free parameter of the approach. It is the same for all low–lying mesons.

The multiplicities of the vector meson production, for example, such as $K^{*\pm}$ and $\rho^{\pm}$ we define as

$$
N_{K^{*+}}(\vec{q}, T) = 3 \times \frac{C_M}{(M_{K^*}F_K)^{3/2}} \int \frac{d^3p}{(2\pi)^3} \frac{1}{e^{-\nu(T)} + |p - \vec{q}|/T + 1} \frac{1}{\sqrt{p^2 + m_{K^*}^2/T} + 1},
$$

$$
N_{K^{*-}}(\vec{q}, T) = 3 \times \frac{C_M}{(M_{K^*}F_K)^{3/2}} \int \frac{d^3p}{(2\pi)^3} \frac{1}{e^{-\nu(T)} + |p - \vec{q}|/T + 1} \frac{1}{\sqrt{p^2 + m_{K^*}^2/T} + 1},
$$

$$
N_{\rho^{\pm}}(\vec{q}, T) = 3 \times \frac{C_M}{(M_\rho F_\rho)^{3/2}} \int \frac{d^3p}{(2\pi)^3} \frac{1}{e^{-\nu(T)} + |p - \vec{q}|/T + 1} \frac{1}{e^{\nu(T)} + p/T + 1},
$$

where $M_{K^*} = 892$ MeV and $M_\rho = 770$ MeV are the masses of the $K^*$ and $\rho$ mesons, respectively [10].

1The numbers of quarks ($q, s$) and anti–quarks ($\bar{q}, \bar{s}$) and the coefficients of proportionality are free parameters of a simple coalescence model. Therefore, a simple coalescence model contains seven free parameters. Five of them can be fixed from experimental data [2].
In the case of baryons and antibaryons we suggest to define the multiplicities by using the diquark–quark picture of baryons and antibaryons. For example, the multiplicities of the proton ($p$) and antiproton ($\bar{p}$) we write in the form

$$N_p(\vec{q}, T) = \frac{3!}{3!} \times \frac{C_B}{(M_p F_\pi)^{3/2}} \times \int \frac{d^3 p}{(2\pi)^3} \frac{1}{(e^{-\nu(T)} + p/T + 1)^2} \frac{1}{e^{-\nu(T)} + |\vec{p} - \vec{q}|/T + 1},$$

$$N_{\bar{p}}(\vec{q}, T) = \frac{3!}{3!} \times \frac{C_{\bar{B}}}{(M_{\bar{p}} F_{\bar{\pi}})^{3/2}} \times \int \frac{d^3 p}{(2\pi)^3} \frac{1}{(e^{\nu(T)} + p/T + 1)^2} \frac{1}{e^{\nu(T)} + |\vec{p} - \vec{q}|/T + 1},$$

where a momentum $\vec{p}$ has a meaning of a relative momentum of three–quark (three–anti–quark) system, $M_p = 940$ MeV is the mass of the proton and antiproton. As well as in the meson case the main contribution to the momentum integrals comes from the momenta of order $p \sim T$ providing a coalescence of three quarks (three anti–quarks) into baryons (anti–baryons) at the momenta of order $p \sim T$. That is again of order of transversal momenta of the produced hadrons, $q_{\perp} \sim (2 \div 3) T$, coupled in the center of mass frame of heavy–ion collisions. The factor $3!$ in the numerator is related to the quark colour degrees of freedom and defined by $\varepsilon_{ijk} \varepsilon^{ijk} = 3!$, where $i, j$ and $k$ are colour indices and run over $i = 1, 2, 3$ each. In turn, in the denominator the factor $3!$ takes into account the identity of three light quarks ($qqq$) and three antiquarks ($\bar{q}\bar{q}\bar{q}$). In the isotopical limit we do not distinguish $u$ and $d$ quarks as well as $\bar{u}$ and $\bar{d}$ antiquarks. The dimensionless parameters $C_B$ and $C_{\bar{B}}$ are free parameters of the approach. Each of them is equal for all components of octets of baryons and antibaryons, respectively, but $C_B \neq C_{\bar{B}}$.

The paper is organized as follows. In Sect. 2 we calculate the theoretical values of multiplicities of the hadron production from the thermalized QGP phase. The theoretical predictions and experimental data are adduced in Table 1. In the Conclusion we discuss the obtained results. A possible estimate of the absolute values of our input parameters is discussed through the application of our approach to the calculation of the number of baryons and antibaryons relative to the number of photons at the early stage of the evolution of the Universe assuming that this evolution goes through the intermediate thermalized QGP phase.

2 Multiplicities of hadron production from the thermalized QGP phase

Now let us proceed to the evaluation of multiplicities of hadron production from the thermalized QGP phase of QCD. The theoretical predictions for the different ratios of hadron multiplicities we compare with experimental data adduced in Table I of Ref. [11]. These are the data of various experiments given by NA44, NA49, NA50 and WA97 Collaborations on Pb+Pb collisions at 158 GeV/nucleon. Also we compare our results with the
experimental data obtained by NA35 Collaboration on S+S collisions and NA38 Collaboration on O+U and S+U collisions at 200 GeV/nucleon. From Table 1 of this paper one can see that in the whole the experimental data for the hadron production are obtained for rapidities ranging over the region $2.3 \leq y \leq 4.1$. The relation between a 3-momentum $q$ and a rapidity $y$ reads

$$q = \sqrt{M^2 \sin^2 y + \vec{q}_\perp^2 c \cos^2 y} \geq M \sin y,$$  \hspace{1cm} (2.1)

where $M$ and $\vec{q}_\perp$ are the mass and the transversal momentum of the produced hadron. For rapidities $y \in [2.3, 4.1]$ we get

$$q \geq (5.0 \div 30.2) M.$$  \hspace{1cm} (2.2)

Thus, for $K$ mesons and hadrons heavier than $K$ mesons typical momenta are of order of 2.5 GeV and greater. This gives a possibility to investigate the momentum integrals defining multiplicities of the hadron production at $q \to \infty$. As has been shown in Ref. [1] the ratios of the multiplicities $R_{K^+ + K^-}(q, T) = N_{K^+ + K^-}(\vec{q}, T)/N_{K^-}(\vec{q}, T)$ and $R_{K^+ + K^-}(q, T) = N_{K^+ + K^-}(\vec{q}, T)/N_{\pi^+}(\vec{q}, T)$ are smooth functions of $q$, wobbling slightly around the asymptotic values obtained at $q \to \infty$, and describe good the experimental data at $T = 175$ MeV. Below the theoretical results on the multiplicities of the hadron production we would compare with experimental data at $T = 175$ MeV. In this case it is obvious that the typical momenta of the momentum integrals defining the multiplicities of hadron production are of order $p \sim T$. Therefore, in our approach the momenta at which quarks and anti–quarks coalesce into hadrons should be of order $p \sim T$. This agrees with the order of transversal momenta of hadrons, $q_\perp \sim 2 \div 3 T$, coupled in the center of mass frame of heavy–ion collisions.

Hence, for rapidities ranging over the region $2.3 \leq y \leq 4.1$ the typical total 3–momenta of hadrons are greater than 1 GeV. Thereby, multiplicities of the hadron production can be calculated in the asymptotic regime at $q \gg T$.

We would like to accentuate that the relative momenta at which quarks and anti–quarks coalesce into hadrons are described effectively by the momentum of integration $\vec{p}$.

At $q \gg T$ the multiplicities of the hadron production defined by Eqs. (1.7)–(1.9) can be represented in the following form

\begin{align*}
N_{\pi^+}(\vec{q}, T) & = N_{\pi^-}(\vec{q}, T) = N_{\pi^0}(\vec{q}, T) = \frac{3 C_M}{(M_{\pi} F_{\pi})^{3/2}} e^{-q/T} I_{\pi}(T), \\
N_{K^+}(\vec{q}, T) & = N_{K^-}(\vec{q}, T) = \frac{3 C_M}{(M_{K} F_{K})^{3/2}} e^{-q/T} e^{+\nu(T)} I_{K}(T), \\
N_{K^-}(\vec{q}, T) & = \frac{1}{2} N_{\pi^0}(\vec{q}, T) + \frac{1}{2} N_{K^0}(\vec{q}, T) = \frac{1}{2} \frac{3 C_M}{(M_{K} F_{K})^{3/2}} e^{-q/T} e^{+\nu(T)} I_{K^0}(T), \\
N_{K_S^0}(\vec{q}, T) & = \frac{3 C_M}{(M_{K_S} F_{K_S})^{3/2}} e^{-q/T} I_{K_S^0}(T), \\
N_{\eta}(\vec{q}, T) & = \sin^2 \theta \frac{3 C_M}{(M_{\eta} F_{\eta})^{3/2}} e^{-q/T} I_{\eta}(T) + \cos^2 \theta \frac{3 C_M}{(M_{S} F_{S})^{3/2}} e^{-q/T} I_{\eta}(T), \\
N_{\phi}(\vec{q}, T) & = \frac{3 C_M}{(M_{\phi} F_{\phi})^{3/2}} e^{-q/T} I_{\phi}(T),
\end{align*}
\[ N_p(\vec{q}, T) = \frac{C_B}{(M_p F_\pi)^{3/2}} e^{-q/T} e^{+\nu(T)} I_p(T), \]
\[ N_\Lambda(\vec{q}, T) = \frac{3 C_B}{(M_\Lambda F_K)^{3/2}} e^{-q/T} e^{+2 \nu(T)} I_\Lambda(T), \]
\[ N_\Xi(\vec{q}, T) = \frac{3 C_B}{(M_\Xi F_S)^{3/2}} e^{-q/T} e^{+\nu(T)} I_\Xi(T), \]
\[ N_\Omega(\vec{q}, T) = \frac{C_B}{(M_\Omega F_S)^{3/2}} e^{-q/T} I_\Omega(T), \]
\[ N_\rho(\vec{q}, T) = \frac{C_B}{(M_\rho F_\pi)^{3/2}} e^{-q/T} e^{-\nu(T)} I_\rho(T), \]
\[ N_\Lambda(\vec{q}, T) = \frac{3 C_B}{(M_\Lambda F_K)^{3/2}} e^{-q/T} e^{-2 \nu(T)} I_\Lambda(T), \]
\[ N_\Xi(\vec{q}, T) = \frac{3 C_B}{(M_\Xi F_S)^{3/2}} e^{-q/T} e^{-\nu(T)} I_\Xi(T), \]
\[ N_\Omega(\vec{q}, T) = \frac{C_B}{(M_\Omega F_S)^{3/2}} e^{-q/T} I_\Omega(T), \]
\[ (2.3) \]

where the structure functions \( I_i(T) (i = \pi, K, \bar{K}, \ldots) \) are defined by

\[
I_\pi(T) = e^{+\nu(T)} \int_0^{\infty} \frac{dp}{4\pi^2} \frac{p^2}{e^{\nu(T)} + p/T + 1} + e^{-\nu(T)} \int_0^{\infty} \frac{dp}{4\pi^2} \frac{p^2}{e^{-\nu(T)} + p/T + 1} = T^3 \frac{3}{4\pi^2} \]
\[
I_K(T) = \int_0^{\infty} \frac{dp}{4\pi^2} \frac{p^2}{e^{\nu(T)} + p/T + 1} + e^{-\nu(T)} \int_0^{\infty} \frac{dp}{4\pi^2} \frac{p^2}{e^{-\nu(T)} + p/T + 1} = T^3 \frac{2.924}{4\pi^2} \]
\[
I_\bar{K}(T) = \int_0^{\infty} \frac{dp}{4\pi^2} \frac{p^2}{e^{+\nu(T)} + p/T + 1} + e^{-\nu(T)} \int_0^{\infty} \frac{dp}{4\pi^2} \frac{p^2}{e^{-\nu(T)} + p/T + 1} = T^3 \frac{3.463}{4\pi^2} \]
\[
I_\eta(T) = I_\phi(T) = \int_0^{\infty} \frac{dp}{2\pi^2} \frac{p^2}{e^{\nu(T)} + p/T + 1} = \frac{m_\pi^2}{2\pi^2} \int_0^{\infty} \frac{dx^2}{e(m_\pi/T)\sqrt{1 + x^2 + 1}} = 3.522 \frac{m_\pi^2}{2\pi^2} \]
\[
I_\rho(T) = \int_0^{\infty} \frac{dp}{2\pi^2} \frac{1}{e^{-\nu(T)} + p/T + 1} = T^3 \frac{1}{2\pi^2} \int_0^{\infty} \frac{dx^2}{e^{-\nu(T)} + x + 1} = 0.253 \frac{T^3}{2\pi^2} \]
\[
I_\Lambda(T) = e^{-\nu(T)} \int_0^{\infty} \frac{dp}{4\pi^2} \frac{p^2}{e^{\nu(T)} + p/T + 1} = e^{-\nu(T)} \int_0^{\infty} \frac{dx^2}{e^{\nu(T)} + x + 1} = 1 \frac{2}{\pi^2} \]

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\[+ e^{-2 \nu(T)} \int_0^\infty \frac{dp}{4\pi^2} \left( \frac{p^2}{(e^{-\nu(T)} + p/T + 1)^2} \right) =
\]
\[= \frac{m_s^3}{4\pi^2} \left[ e^{-\nu(T)} \int_0^\infty \frac{dx}{e(1 + x^2)} \right] = \frac{1}{e^{-\nu(T)} + (m_s/T) x + 1}
\]
\[+ e^{-2 \nu(T)} \int_0^\infty \frac{dx}{(e^{-\nu(T)} + (m_s/T) x + 1)^2} = 0.582 \frac{m_s^3}{4\pi^2}.
\]
\[I_{\dot{z}}(T) = \frac{dp}{4\pi^2} \frac{p^2}{(e\nu(T) + p/T + 1)^2}
\]
\[+ e^{-\nu(T)} \int_0^\infty \frac{dx}{e\nu(T) + p/T + 1} = \frac{m_s^3}{4\pi^2} \left[ e^{-\nu(T)} \int_0^\infty \frac{dx}{e(1 + x^2)} \right] = \frac{1}{e^{-\nu(T)} + (m_s/T) x + 1}
\]
\[= 0.528 \frac{m_s^3}{4\pi^2}.
\]
\[I_{\dot{\theta}}(T) = \frac{dp}{2\pi^2} \frac{p^2}{(e\nu(T) + p/T + 1)^2} = \frac{T^3}{2\pi^2} \int_0^\infty \frac{dx}{(e\nu(T) + x + 1)^2} = 0.097 \frac{T^3}{2\pi^2},
\]
\[\int_0^\infty \frac{dx}{(e\nu(T) + x + 1)^2} = 0.097 \frac{T^3}{2\pi^2},
\]
\[I_{\dot{\ell}}(T) = \frac{dp}{2\pi^2} \frac{p^2}{(e\nu(T) + p/T + 1)^2} = \frac{T^3}{2\pi^2} \int_0^\infty \frac{dx}{(e\nu(T) + x + 1)^2} = 0.097 \frac{T^3}{2\pi^2},
\]
\[\int_0^\infty \frac{dx}{(e\nu(T) + x + 1)^2} = 0.097 \frac{T^3}{2\pi^2},
\]
\[I_{\dot{\lambda}}(T) = e^{+\nu(T)} \int_0^\infty \frac{dp}{4\pi^2} \frac{p^2}{e\nu(T) + p/T + 1} \frac{1}{e^{+\nu(T)} + (m_s/T) x + 1}
\]
\[+ e^{+2 \nu(T)} \int_0^\infty \frac{dx}{(e^{+\nu(T)} + (m_s/T) x + 1)^2} = \frac{m_s^3}{4\pi^2} \left[ e^{+\nu(T)} \int_0^\infty \frac{dx}{e(1 + x^2)} \right] = 0.588 \frac{m_s^3}{4\pi^2}.
\]
\[I_{\dot{z}}(T) = \frac{dp}{4\pi^2} \frac{p^2}{(e\nu(T) + p/T + 1)^2}
\]
\[ + e^{+ \nu(T)} \int_0^\infty dp \frac{p^2}{4\pi^2 e^{p^2/m_s^2/T}} \frac{1}{e^{p/\nu(T) + p/T} + 1} = \]
\[ = \frac{m_s^3}{4\pi^2} \left[ \int_0^\infty dx \frac{x^2}{(e^{m_s/T})^2 + x^2 + 1} \right] = e^{+ \nu(T)} \int_0^\infty dx \frac{x^2}{e^{m_s/T}(1 + x^2 + 1)} \times \frac{1}{e^{\nu(T)} + \frac{m_s^3}{4\pi^2}} = 0.558 \frac{m_s^3}{4\pi^2}. \]

\[ I_{\Omega}(T) = I_{\Omega}(T). \] (2.4)

The numerical values of the integrals are obtained at \( m_s = 135 \text{ MeV} \) and \( T = 175 \text{ MeV} \).

The theoretical ratios of multiplicities of the hadron production which we compare with measured experimentally we define as follows

\[ R_{K^+K^-}(q, T) = \frac{N_{K^+}(q, T)}{N_{K^-}(q, T)} = e^{2\nu(T)} \frac{I_{K}(T)}{I_{\pi}(T)} = 1.520, \]
\[ R_{K^+\pi^+}(q, T) = \frac{N_{K^+}(q, T)}{N_{\pi^+}(q, T)} = \left( \frac{M_{\pi} F_{\pi}}{M_{K} F_{K}} \right)^{3/2} e^{\nu(T)} \frac{I_{K}(T)}{I_{\pi}(T)} = 0.139, \]
\[ R_{K^-\pi^-}(q, T) = \frac{N_{K^-}(q, T)}{N_{\pi^-}(q, T)} = \left( \frac{M_{\pi} F_{\pi}}{M_{K} F_{K}} \right)^{3/2} e^{-\nu(T)} \frac{I_{K}(T)}{I_{\pi}(T)} = 0.090, \]
\[ R_{K_S^0\pi^-}(q, T) = \frac{N_{K_S^0}(q, T)}{N_{\pi^-}(q, T)} = \frac{1}{2} \left( \frac{M_{\pi} F_{\pi}}{M_{K} F_{K}} \right)^{3/2} e^{\nu(T)} \frac{I_{K_S^0}(T)}{I_{\pi^-}(T)} = 0.113, \]
\[ R_{\Xi\Lambda}(q, T) = \frac{N_{\Xi}(q, T)}{N_{\Lambda}(q, T)} = \left( \frac{M_{\Xi} F_{\Xi}}{M_{\Lambda} F_{\Lambda}} \right)^{3/2} e^{-\nu(T)} \frac{I_{\Xi}(T)}{I_{\Lambda}(T)} = 0.108, \]
\[ R_{0\Xi}(q, T) = \frac{N_{0\Xi}(q, T)}{N_{\Xi}(q, T)} = \frac{1}{3} \left( \frac{M_{\Xi}}{M_{0}} \right)^{3/2} e^{-\nu(T)} \frac{I_{0\Xi}(T)}{I_{\Xi}(T)} = 0.166, \]
\[ R_{\Lambda\Phi}(q, T) = \frac{N_{\Lambda}(q, T)}{N_{\Phi}(q, T)} = 3 \left( \frac{M_{\Phi} F_{\Phi}}{M_{\Lambda} F_{\Lambda}} \right)^{3/2} e^{-\nu(T)} \frac{I_{\Lambda}(T)}{I_{\Phi}(T)} = 2.081, \]
\[ R_{\Xi\Lambda}(q, T) = \frac{N_{\Xi}(q, T)}{N_{\Lambda}(q, T)} = \left( \frac{M_{\Xi} F_{\Xi}}{M_{\Lambda} F_{\Lambda}} \right)^{3/2} e^{\nu(T)} \frac{I_{\Xi}(T)}{I_{\Lambda}(T)} = 0.173, \]
\[ R_{0\Xi}(q, T) = \frac{N_{0\Xi}(q, T)}{N_{\Xi}(q, T)} = \frac{1}{3} \left( \frac{M_{\Xi}}{M_{0}} \right)^{3/2} e^{\nu(T)} \frac{I_{0\Xi}(T)}{I_{\Xi}(T)} = 0.282, \]
\[ R_{0\Omega}(q, T) = \frac{N_{0\Omega}(q, T)}{N_{\Omega}(q, T)} = \frac{C_{\Omega}}{C_{0}} = \frac{R_{0\Omega}}{C_{0}} = 0.46 \pm 0.15, \]
\[ R_{\Lambda\Lambda}(q, T) = \frac{N_{\Lambda}(q, T)}{N_{\Lambda}(q, T)} = \frac{C_{\Phi}}{C_{0}} \times e^{-4\nu(T)} \frac{I_{\Lambda}(T)}{I_{\Lambda}(T)} = 0.168 \pm 0.055, \]
\[ R_{\Xi\Xi}(q, T) = \frac{N_{\Xi}(q, T)}{N_{\Xi}(q, T)} = \frac{C_{\Phi}}{C_{0}} \times e^{-2\nu(T)} \frac{I_{\Xi}(T)}{I_{\Xi}(T)} = 0.270 \pm 0.088, \]
\[ R_{\eta\eta^0}(q, T) = \frac{N_{\eta}(q, T)}{N_{\eta^0}(q, T)} = \sin^2 \hat{\theta} \left( \frac{M_{\eta}}{M_{\eta^0}} \right)^{3/2} + \cos^2 \hat{\theta} \left( \frac{M_{\pi} F_{\pi}}{M_{\eta} F_{\eta}} \right)^{3/2} \frac{I_{\eta}(T)}{I_{\eta^0}(T)} = 0.088, \]
by Bramon, Escribano and Scadron [13] gives mixing angle. Recent analysis of the value of the octet–singlet mixing angle carried out

\[ \eta \]

The multiplicity of the \( \eta \) meson phenomenology [10,12,13] gives the following quark structure of the \( \eta \):

\[ \eta(550) = (q\bar{q}) \sin \theta + (s\bar{s}) \cos \theta, \]  

where \( \theta = \theta_0 - \theta_P \) with \( \theta_0 = 35.264^\circ \), the ideal mixing angle, and \( \theta_P \), the octet–singlet mixing angle. Recent analysis of the value of the octet–singlet mixing angle carried out by Bramon, Escribano and Scadron [13] gives \( \theta_P = -16.9 \pm 1.7^\circ \). For the \( \phi(1020) \) meson we have supposed the \( s\bar{s} \) quark structure [10,12].

### 3 Conclusion

The theoretical and experimental values of the ratios of hadron production are adduced in Table 1. From Table 1 one can see a good agreement between presently available set of hadron ratios measured for various experiments given by NA44, NA49, NA50 and WA97 Collaborations on Pb+Pb collisions at 158 GeV/nucleon, NA35 Collaboration on S+S collisions and NA38 Collaboration on O+U and S+U collisions at 200 GeV/nucleon and theoretical predictions for the ratios of multiplicities of hadron production from the thermalized QGP phase at a temperature \( T = 175 \text{ MeV} \). Save the ratio \( \Lambda/\bar{p}, (\Lambda/\bar{p})_{\text{th}} = 2.081 \) and \( (\Lambda/\bar{p})_{\text{exp}} = 3 \pm 1 \), the theoretical results agree with the experimental ones with accuracy better than 18%.

In our approach multiplicities of hadron production are defined by momentum integrals on quark (antiquark) distribution functions in accordance with the phenomenological quark structure of the hadron. For the analysis of the multiplicities of the baryon
and antibaryon production in terms of the quark and antiquark distribution functions we have followed the diquark–quark picture for baryons and antibaryons. This has allowed to describe multiplicities of the baryon, antibaryon and meson production on the same footing.

For the explanation of experimental data on the hadron production in ultrarelativistic heavy–ion collisions we have used three input parameters $C_B/C_B$, $C_M/C_B$ and $F_S$. These parameters are related to the spatial volumes of hadronization of the quarks and antiquarks from the thermalized QGP phase. The first two parameters have been fixed from the experimental data on the ratios $(\bar{\Omega}/\Omega)_{\text{exp}} = 0.46 \pm 0.15$, $(\Lambda/K_S^0)_{\text{exp}} = 0.65 \pm 0.11$. This gives $C_B/C_B = 0.46 \pm 0.15$ and $C_M/C_B = 0.23 \pm 0.04$. In turn, the value of the parameter $F_S = 3.5 F_\pi = 458.5 \text{MeV}$ is a result of a smooth fit of the ratios of hadrons containing the $ss$ and $\bar{s}s$ components in the quark structure. In the bulk our approach to the hadronization from the QGP phase of QCD has succeeded in describing 21 experimental data on ultrarelativistic heavy–ion collisions.

Unlike other approaches [2–4] the ratio of the $\bar{\Omega}$ and $\Omega$ baryon production is an input parameter in our model $C_B/C_B$. The ratio $\bar{\Omega}/\Omega$ does not depend in our approach on both the momenta of baryons and the temperature. The former is due to the zero–value of the $s$–quark chemical potential, $\mu_s = \mu_{\bar{s}} = 0$. As a result the ratio $\bar{\Omega}/\Omega$ can be only fitted in our approach. By fitting the ratio $\bar{\Omega}/\Omega$ from experimental data and applying this value to the description of other ratios of the baryon and antibaryon production for the thermalized QGP phase we have found a good agreement with experimental data. This confirms a self–consistency of the approach.

In turn, the distinction between the parameters $C_B$ and $C_B$ can be related to a well–known factor of the baryon–antibaryon asymmetry in the Universe which one could put phenomenologically at the early stage of the evolution of the Universe whether the baryon synthesis in it goes through the intermediate QGP phase. Indeed, as has been stated by Börner [29]: *Within the standard big–bang model, however, there seems to be little chance of achieving a physical separation of baryon and antibaryon phases in an initially baryon–symmetric cosmological model. If the baryon number was exactly conserved – as it assumed to be the standard model – the small asymmetry necessary for our existence must be postulated initially. Grand unified theories offer the possibility of creating this small asymmetry from physical processes.*

In our approach the baryon–antibaryon asymmetry at the hadronic level can be realized phenomenologically in terms of a different rate of hadronization of baryons and antibaryons caused by the input parameter $C_B/C_B = 0.46 \pm 0.15$ fixed through the experimental data on the ratio $\bar{\Omega}/\Omega$ production in ultrarelativistic heavy–ion collisions. For the total number of antibaryons $N_B$ relative to the total number of baryons $N_B$ produced for the baryon–antibaryon synthesis at the early stage of the evolution of the Universe at a temperature $T = 175 \text{MeV}$ [30] and gone through the intermediate thermalized QGP phase we predict

$$\frac{N_B}{N_B} = 0.41 \times \frac{C_B}{C_B} = 0.19 \pm 0.06.$$  (3.1)

This result can be supported by a trivial estimate in the approximation of the equilibrium
baryon and anti–baryon gases. In such an approximation the ratio \( N_{\bar{B}}/N_B \) is defined by

\[
\frac{N_{\bar{B}}}{N_B} = e^{-2\mu_B(T)/T} = e^{-6\mu(T)/T} = 0.17,
\]

where a chemical potential \( \mu(T) \) is given by Eq. (1.4) and \( T = 175 \text{ MeV} \).

Thus, at the early stage of the evolution of the Universe the number of antibaryons should be of order of magnitude less compared with the number of baryons, \( N_{\bar{B}} \sim 0.2 N_B \).

According to Börner [31] it is more than enough for the existence of the life in the Universe. Recall that the standard approach [31] predicts for every \( 10^9 \) antibaryons only \( (10^9 + 1) \) baryons. As has been stated by Börner: *It is to that one part in \( 10^9 \) excess of ordinary matter that we owe our existence!* [31].

For the early Universe the total number of baryons and antibaryons was roughly equal to the number of photons \( N_{ph} \) [31]:

\[
\frac{N_{\bar{B}} + N_B}{N_{ph}} \simeq 1.
\]

Since the density of photons is equal to [32]

\[
\frac{N_{ph}}{V} = \frac{2.404}{\pi^2} T^3,
\]

where \( V \) is the volume of the early Universe, and the density of the total number of baryons and antibaryons \( N_{\bar{B}} + N_B \) calculated in our approach at \( T = 175 \text{ MeV} \) amounts to

\[
\frac{N_{\bar{B}} + N_B}{V} = C_B \times 0.442 \times T^3,
\]

we can estimate the numerical value of the parameter \( C_B \):

\[
C_B \simeq 0.55 \pm 0.08.
\]

This gives the estimate of other input parameters \( C_{\bar{B}} \) and \( C_M \):

\[
C_{\bar{B}} \simeq 0.25 \pm 0.08, \\
C_M \simeq 0.13 \pm 0.03.
\]

The analysis of the influence of the input parameter \( C_B/C_B = 0.46 \pm 0.15 \) on the evolution of the baryon–antibaryon asymmetry in the Universe from the early Universe up to the present epoch and the formation of a dark and strange matter in the Universe [10] we are planning to carry out in our forthcoming publications.

Now we would like to discuss in more details our approach when compared with a simple coalescence model [2]. The main distinction of our approach from a simple coalescence model is in the correlation between quarks and anti–quarks coalescing into hadrons. In fact, in a simple coalescence model quarks and anti–quarks are uncorrelated [2]. This has allowed to introduce separately the number of light quarks \( q \) and light anti–quarks \( \bar{q} \) and the number of strange quarks \( s \) and strange anti–quarks \( \bar{s} \) [2]. Moreover,
this has turned out to be of use in order to hide the dependence of quark and anti–quark distribution functions on a temperature \( T \) and a chemical potential \( \mu(T) \) in the number of light quarks \( q \) and light anti–quarks \( \bar{q} \). A non–vanishing chemical potential of strange quarks \( \mu_s(T) \) is hidden in the number of strange quarks \( s \) and strange anti–quarks \( \bar{s} \). Then, the calculation of multiplicities of hadrons produced from the QGP phase in a simple coalescence model resembles a quark counting. In fact, multiplicities of hadrons are proportional to products of the number of quarks \( q, s \) and anti–quarks \( \bar{q}, \bar{s} \) in accord the naive quark structure of hadrons. Since quarks and anti–quarks do not correlate, so that the multiplicities of hadrons turn out to be independent on the momenta of hadrons. Then, the numbers of quarks \( q, s \) and anti–quarks \( \bar{q}, \bar{s} \) and the coefficients of proportionality, the coalescence coefficients \( C_p, C_{\Lambda}, C_{\Xi}, C_{\Omega} \) and \( C_{\bar{p}}, C_{\bar{\Lambda}}, C_{\bar{\Xi}}, C_{\bar{\Omega}} \), are free parameters of a simple coalescence model. Therefore, a total number free parameters appearing in a simple coalescence model for the description of baryon and anti–baryon production is equal to twelve. By the assumption \( C_p/C_{\bar{p}} = C_{\Lambda}/C_{\bar{\Lambda}} = C_{\Xi}/C_{\bar{\Xi}} = C_{\Omega}/C_{\bar{\Omega}} = 1 \) the number free parameters has been reduced up to five \( (q, s, \bar{q}, \bar{s}, C) \), where \( C \) is a common for baryons and anti–baryons coalescence coefficient. Two of these free parameters have been fixed from the fit of experimental data on the baryon and anti–baryon production: \( \bar{q}/q = 0.41 \pm 0.02 \) and \( \bar{s}/s = 0.75 \pm 0.06 \) [2]. Thus, there are three free parameters left in a simple coalescence model applied to the description of baryon and anti–baryon production from the QGP phase. It is also important to note that a simple coalescence model [2] explains only multiplicities of baryon and anti–baryon production. In fact, save the ratio of multiplicities of the \( K^+ \) and \( K^- \) mesons none other multiplicities of pseudoscalar and vector mesons have been predicted within a simple coalescence model [2]. Therefore, it is not completely clear how many free parameters would be added in a simple coalescence model for description of multiplicities of pseudoscalar and vector meson production.

In our approach, where quarks and anti–quarks coalescing into hadrons are correlated, we have six free parameters \( T, \mu(T), C_M, C_B, C_{\bar{B}} \) and \( F_S \). Five of these parameters \( T = 175 \text{ MeV}, \mu(T) \) given by Eq. (1.4), \( F_S = 3.5 F_\pi = 458.5 \text{ MeV}, C_M/C_B = 0.23 \pm 0.04 \) and \( C_{\bar{B}}/C_B = 0.46 \pm 0.15 \) are fixed. Therefore, only there is one free parameter left in the approach. Thus, if to take into account that within our approach we have described not only multiplicities of baryon and anti–baryon production from the QGP but also multiplicities of pseudoscalar and vector meson production, all together 21 experimental ratios, our approach to the thermalized QGP looks much more preferable with respect to a simple coalescence model. If to remind that in our approach due to correlations between quarks and anti–quarks we are able to follow a dependence of multiplicities of hadron production on the hadronic momenta, so an advantage of our approach with respect to a simple coalescence model becomes obvious.

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Table 1. The theoretical ratios of multiplicities of the hadron production are compared with the experimental data obtained by NA44, NA49, NA50 and WA97 Collaborations on Pb+Pb collisions at 158 GeV/nucleon, NA35 Collaboration on S+S collisions and NA38 Collaboration on O+U and S+U collisions at 200 GeV/nucleon. The theoretical multiplicities are calculated at the temperature \( T = 175 \text{ MeV} \).

| N  | Ratio       | Model | Data       | Coll. | Rapidity | Ref. |
|----|-------------|-------|------------|-------|----------|------|
| 1  | \( \bar{p}/p \) | 0.097(32) | 0.055(10) | NA44  | 2.3–2.9  | [14] |
|    |             | 0.097(32) | 0.085(8) | NA49  | 2.5–3.3  | [15] |
| 2  | \( \bar{\Lambda}/\Lambda \) | 0.168(55) | 0.128(12) | WA97  | 2.4–3.4  | [16] |
| 3  | \( \bar{\Xi}/\Xi \) | 0.270(88) | 0.227(33) | NA49  | 3.1–3.85 | [17] |
|    |             | 0.270(88) | 0.266(28) | WA97  | 2.4–3.4  | [16] |
| 4  | \( \bar{\eta}/\eta \) | fit | 0.46(15) | WA97  | 2.4–3.4  | [16] |
| 5  | \( \Xi/\Lambda \) | 0.108 | 0.127(11) | NA49  | 3.1–3.85 | [17] |
|    |             | 0.108 | 0.093(7)  | WA97  | 2.4–3.4  | [16] |
| 6  | \( \Omega/\Xi \) | 0.166 | 0.195(28) | WA97  | 2.4–3.4  | [16] |
| 7  | \( \bar{\Xi}/\bar{\Lambda} \) | 0.173 | 0.180(39) | NA49  | 3.1–3.85 | [17] |
|    |             | 0.173 | 0.195(23) | WA97  | 2.4–3.4  | [16] |
| 8  | \( \bar{\Lambda}/\bar{p} \) | 2.081 | 3(1) | NA49  | 3.1–3.85 | [18] |
| 9  | \( \bar{\Omega}/\Xi \) | 0.282 | 0.27(6)  | WA97  | 2.4–3.4  | [19] |
| 10 | \( K^+/K^- \) | 1.520 | 1.85(9)  | NA44  | 2.4–3.5  | [14] |
|    |             | 1.520 | 1.8(1)   | NA49  | all      | [20] |
| 11 | \( K^+/#pi^+ \) | 0.139 | 0.137(8) | NA35  | all      | [21] |
| 12 | \( K^-/#pi^- \) | 0.090 | 0.076(5) | NA35  | all      | [21] |
| 13 | \( K_S^0/#pi^- \) | 0.113 | 0.125(19) | NA49  | all      | [22] |
| 14 | \( \eta/#pi^0 \) | 0.088 | 0.081(13) | WA98  | 2.3–2.9  | [23] |
| 15 | \( 2\phi/(#pi^+ + #pi^-) \) | \( 7.8 \times 10^{-3} \) | \( 9.1(1.0) \times 10^{-3} \) | NA50  | 2.9–3.9  | [24] |
| 16 | \( \phi/(#rho^0 + #omega^0) \) | 0.103 | \( \approx 0.1 \) | NA38  | 2.8–4.1  | [25,26] |
| 17 | \( \phi/#K_S^0 \) | 0.071 | 0.084(11) | NA49  | all      | [27] |
| 18 | \( \Lambda/#K_S^0 \) | fit | 0.65(11) | WA97  | 2.4–3.4  | [28] |
| 19 | \( p/#K^+ \) | 0.136(23) | 0.102(19) | NA44  | 2.3–2.9  | [14] |
| 20 | \( \bar{p}/K^- \) | 0.020(7) | 0.153(17) | NA49  | 2.5–3.3  | [15,20] |
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