Vol. 29 (1998) *ACTA PHYSICA POLONICA B* No 12

QUARK–GLUON PLASMA*

Stanisław Mrówczyński

Sołtan Institute for Nuclear Studies
Hoża 69, 00-681 Warsaw, Poland
and
Institute of Physics, Pedagogical University
Konopnickiej 15, 25-406 Kielce, Poland
e-mail: mrow@fuw.edu.pl

(Received October 26, 1998)

An elementary introduction to the physics of quark–gluon plasma is given. We start with a sketchy presentation of the Quantum Chromodynamics which is the fundamental theory of strong interactions. The structure of hadrons built up of quarks and gluons is briefly discussed with a special emphasis on the confinement hypothesis. Then, we explain what is the quark–gluon plasma and consider why and when the hadrons can dissolve liberating the quarks and gluons. The heavy-ion collisions at high-energies, which provide a unique opportunity to get a droplet of the quark–gluon plasma in the terrestrial conditions, are described. We also consider the most promising experimental signatures of the quark–gluon plasma produced in nucleus–nucleus collisions. At the end, the perspectives of the quark–gluon plasma studies at the future accelerators are mentioned.

PACS numbers: 12.38.Mh, 25.75.+r

1. Introduction

The quark–gluon plasma is a state of the extremely dense matter with the quarks and gluons being its constituents. Soon after the Big Bang the matter was just in such a phase. When the Universe was expanding and cooling down the quark–gluon plasma turned into hadrons — neutrons and protons, in particular — which further formed the atomic nuclei. The plasma is not directly observed in Nature nowadays, but the astrophysical compact objects as the neutron stars may conceal the quark–gluon nuggets in their dense centers. The most exciting however are the prospects to study the plasma

* Presented at the XXXVIII Cracow School of Theoretical Physics, Zakopane, Poland, June 1–10, 1998.

in the laboratory experiments. A broad research program of the heavy-ion collisions, which provide a unique opportunity to produce the quark–gluon droplets in the terrestrial conditions, is underway. While the question, whether the plasma is already produced with the currently operating heavy-ion accelerators, is right now vigorously debated, there are hardly any doubts that we will have a reliable evidence of the plasma within a few years. Then, there will be completed the accelerators of a new generation: the Relativistic Heavy Ion Collider (RHIC) at the Brookhaven National Laboratory and the Large Hadron Collider (LHC) at CERN.

The aim of this article is to give a very elementary introduction to the physics of quark–gluon plasma. We start with a few words on the Quantum Chromodynamics which is a fundamental theory of strong interactions. Then, the structure of hadrons, which are built up of quarks and gluons, is briefly discussed. A particular emphasis is paid on the confinement hypothesis. In the next section we try to explain why and when the hadrons can dissolve liberating the quarks and gluons from their interiors. Finally, the generation of the quark–gluon plasma in heavy-ion collisions is considered.

There is a huge literature on every topic touched in this article. Instead of citing numerous original papers, we rather recommend two collections of the review articles [1,2]. The more recent progress in the field can be followed due to the proceedings of the regular "Quark Matter" conferences [3–5].

Throughout the article we use the natural units where the velocity of light c, the Planck $\hbar$ and Boltzmann k constants are all equal unity. Then, the mass, momentum and temperature have the dimension of energy, which is most often expressed in MeV. The distances in space or time are either measured in the length units, which are usually given in fm (1 fm $= 10^{-13}$ cm), or in the inverse energy units. One easily recalculates the length into the inverse energy or vice verse keeping in mind that $\hbar c = 197.3$ MeV$\cdot$ fm.

2. Quantum chromodynamics

Quantum Chromodynamics (QCD) strongly resembles the Quantum Electrodynamics (QED). While QED describes the interaction of electric charges — usually electrons and their antiparticles positrons — with the electromagnetic field represented by photons, QCD deals with the quarks and gluons corresponding to, respectively, the electrons and photons. The quarks are, as electrons, massive and curry a specific charge called color, which is however not of one but of three types: red, blue and green. The gluons, which are, as photons, massless, but not neutral. In contrast to photons, they curry color charges being the combinations of the quark ones. The electromagnetic interaction proceeds due to the photon exchanges. Analogously the quarks interact exchanging the gluons. Although the photons being neutral cannot interact directly with each other, there are forces acting between gluons.

QCD has emerged as a theory of quarks and gluons which built up the hadrons *i.e.* strongly interacting particles such as neutrons and protons. However, there is no commonly accepted model of the hadron structure. The difficulty lies in the very nature of the strong interaction — its strength. While the perturbative expansion, where the noninteracting system is treated as a first approximation, appear to be the only effective and universal computational method in the quantum field theory, a large value of the QCD coupling constant excludes applicability of the method for the system of quarks and gluons. However, QCD possesses a remarkable property called the *asymptotic freedom*. The coupling constant α_s effectively depends on the four-momentum Q transferred in the interaction as

$$\alpha_s(Q) = \frac{12\pi}{(33-2N_f)\,\ln Q^2/\Lambda_{\rm QCD}^2}\ , \tag{1}$$

where N_f is the number of the number of the quark flavours (types) and $\Lambda_{\rm QCD}$ is the QCD scale parameter, $\Lambda_{\rm QCD} \cong 200$ MeV. Eq. (1) shows that the coupling constant is small when $Q^2 \gg \Lambda_{\rm QCD}^2$. Therefore, the interactions with a large momentum transfer can be treated in the perturbative way. QCD appears to be indeed very successful in describing the hard processes, such as a production of jets in high energy proton–antiproton collisions, which proceed with high Q.

3. Hadron structure

The description of the soft processes in QCD, which, in particular, control the hadron structure, remains a very serious unresolved problem of the strong interactions. One has to rely on the phenomenological models, the validity of which can be only tested by confrontation against the experimental data. Dealing with the soft QCD one often uses the concept of the constituent quarks which should be distinguished from the current quarks. The latter ones are the fundamental elementary spin 1/2 particles which are present in the QCD lagrangian. The current up (u) and down (d) quarks are light with the mass of a few MeV. The constituent u and d quarks are supposed to be the effective quasiparticles with the (large) masses of about 300 MeV (1/3 of the nucleon) generated by the interaction. The difference between the current and constituent strange (s) quark masses is less dramatic. They are, respectively, about 150 and 450 MeV. In the case of heavy quarks — charm (c), bottom (b) and top (t) — the distinction is no longer valid.

Within the constituent quark model the baryon is described as three bounded quarks. The meson is then the system of quark and antiquark.

In terms of current quarks the hadron is seen as a cloud of quarks and gluons. The baryon is then no longer built of three quarks and the meson of quark and antiquark. Instead, the hadron curries the quantum numbers of, respectively, three quarks or the quark and antiquark. Therefore, the hadron is composed of the valence quarks (quark and antiquark in the case of meson and three quarks for the baryon) and the sea constituted by the quark–antiquark pairs and gluons.

The gluons are believed to glue together the quarks which form the hadron, but a satisfactory theory of the hadron binding is still missing. Such a theory has to explain the *hypothesis of confinement* which is the fundamental element of our understanding of the hadron world. While the electric charges tend to form electrically neutral atoms and molecules, in the case of the chromodynamic interactions there seems to be a strict rule that the color charges occur only within the white configurations. The existence of the separated color objects as quarks and gluons is excluded. They must be confined in the colorless systems such as hadrons. The three quarks forming a baryon curry three fundamental colors which give all together the white object. In the meson case the quark color is complementary to the color of the antiquark. In principle, the confinement hypothesis allows for the existence of not only the meson (quark–antiquark) and baryon (three quark) configurations but for any white one such as a dibaryon which is the six quark system. However, in spite of hard experimental efforts, the reliable evidence for the hadrons different than baryons and mesons is lacking.

There are many phenomenological approaches to the confinement. Let us briefly present here the string model inspired by the Meissner effect which is the expelling of the magnetic field from the superconducting material. The model assumes that the vacuum behaves as a diaelectric medium, where the chromodynamic field cannot propagate but is confined in thin tubes or strings which connect the field sources. In Fig. 1 we show the electric field generated by the two opposite charges which are in the (normal) vacuum (a) and in the diaelectric medium (b). Let us compute the potential which acts between charges in the latter case. Using the Gauss theorem one immediately finds the electric field as $E = q/\sigma$, where q denotes the charge and σ the cross section of the tube. If σ is independent of the distance r between the charges, their potential energy equals

$$V(r) = \frac{q^2}{\sigma}\, r \, . \tag{2}$$

As seen the potential energy grows lineary with r when the charges are put in the diaelectric medium. Keeping in mind the result (2), the confinement hypothesis, which is illustrated in Fig. 2, can be understood as follows. Let us imagine that one tries to burst the meson up separating the quark from

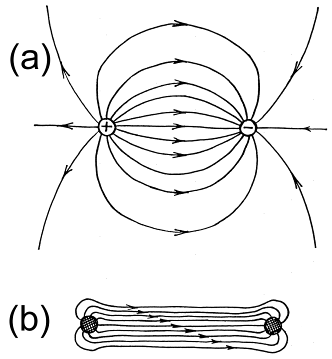

Fig. 1. The electric field lines in the vacuum (a) and diaelectric medium (b)

the antiquark. Stretching the meson requires pumping of the energy to the system. When the energy is sufficient to produce the quark–antiquark pair, the string breaks down and we have two mesons instead of one.

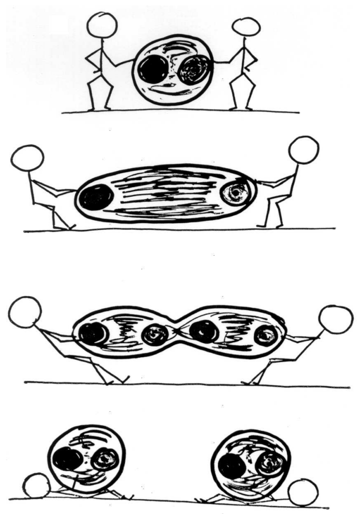

Fig. 2. The confinement

At the end of this section we mention another commonly used model of the hadron structure: the bag model. One assumes here that the vacuum exerts the pressure B on the (colored) quarks and gluons. Then, the hadron is the bag with quarks and gluons similar to the bubble of vapour in the liquid. The bag is usually spherical but the deformations are possible. The model appears to be really successful in describing a very reach mass spectrum of hadrons.

4. Quark–gluon plasma

The quark–gluon plasma is the system of quarks and gluons, which are no longer confined in the hadron interiors, but can propagate in the whole volume occupied by the system. Thus, the quark–gluon plasma resembles the ionized atomic gas with quarks and gluons corresponding to electrons and ions and hadrons being the analogs of atoms. One may wonder whether the existence of quark–gluon plasma is not in conflict with the confinement hypothesis. We note that the plasma is white as a whole. Thus, the color charges are still confined in the colorless system. However, it is still unclear why and when the hadrons can dissolve liberating quarks and gluons. We will consider this issue in the next section. Here we would like to discuss a somewhat unexpected consequence of the asymptotic freedom.

Due to the mentioned difficulties of the soft QCD, the properties of the hadronic matter — either composed of hadrons or of quarks and gluons — are poorly known at the moderate temperatures. The situation changes qualitatively when the system temperature T is much larger than the QCD scale parameter $\Lambda_{\rm QCD}$. In this limit T is the only dimensional parameter which describes the system. In particular, it determines the average momentum transfer in the interaction of quarks and gluons. Therefore, $\langle Q^2 \rangle = c\,T^2$ where c is a dimensionless constant. On the basis of the dimensional argument we expect to achieve the asymptotic freedom regime (smallness of the coupling constant) when $T \gg \Lambda_{\rm QCD}$. Then, the quark–gluon plasma is a weakly interacting gas of massless quarks and gluons.

The detailed calculations performed within the thermal QCD show that the asymptotic freedom regime is indeed obtained in the high temeprature limit. Here we will only add a simple physical argument to the dimensional one used above. The momentum transfer, which enters the formula of the running coupling constant (1), corresponds (due to the Fourier transform) to the distance between the interacting partons. The distance is proportional to the inverse momentum transfer. When the plasma temperature increases, the density of the quark–gluon system grows as T^3 in the high temperature limit. (One can refer here again to the dimensional arguments or simple calculations presented in the next section.) Since the average particle separation in the gas of the density ρ equals $\rho^{-1/3}$, the average inter-particle distance decreases with T as T^{-1}. Consequently, the average coupling constant vanishes when $T \rightarrow \infty$.

5. Deconfinement phase transition

The transformation of the hadron gas into the quark–gluon plasma is called the deconfinement phase transition. There are many independent indications that such a transition indeed takes place in the dense hadronic medium. First of all one should mention the results obtained within the lattice formulation of QCD, where the space continuum is replaced by the discrete points. The Monte Carlo simulations show that there are two phases in the lattice QCD, which are identified with the hadron and quark–gluon phase, respectively. Here we would like to discuss simple arguments in favour of the existence of the quark–gluon plasma.

We start with the observation that hadrons are not point-like objects but their size is finite. The hadron radius is about 1 fm. So, let us consider the hadron gas which is so dense that the average separation of hadrons is about 1 fm. There is no reason to think that the confining potential still operates in such a medium at the distances which are significantly larger than 1 fm. When the quark and antiquark forming a meson are pulled way from each other there is no vacuum, which expels the chromodynamic field, but there are other hadrons all around. Therefore, the confining potential is expected to be screened at the distances comparable to the average inter-particle separation.

There are two ways, as illustrated in Fig. 3, to produce the dense hadronic matter. The first one is evident — squeezing of the nuclear matter. Since the baryon number, which is curried by neutrons and protons, is conserved, the nucleons cannot disappear and will start to overlap when the average inter-particle distance is smaller than the nucleon radius $r_N \cong 1$ fm. The inter-nucleon separation is smaller than r_N at the densities exceeding $\rho = r_N^{-3} \cong 1\ \mathrm{fm}^{-3} \cong 6\,\rho_0$, where $\rho_0 = 0.16\ \mathrm{fm}^{-3}$ is the so-called saturation or normal nuclear density being approximately equal the nucleon density in the nucleus center.

The second method to produce the dense hadronic matter is to heat up the nuclear matter or the hadron gas. The point is that in the contrast to the baryon or lepton number, which are the conserved quantities, the particle number is not. Therefore, when the gas temperature (measured in the energy units) becomes comparable to the particle mass, further heating leads not only to the increase of the average particle kinetic energy but to the growth of the average particle number. Of course, the particle number growth cannot violate the conservation laws. Therefore, the particles, which curry a conserved charge, must be produced in the particle–antiparticle pairs. For example, to keep the baryon number of the system fixed we can add to the system only pairs of the baryon and antibaryon. The essentially neutral particles such as γ or π^0, which do not curry conserved charges, can be added to the system without restrictions, although the average particle number is controlled by the equilibrium conditions. More specifically, the numbers are determined by the minimum of the system free energy when the system temperature and volume are fixed.

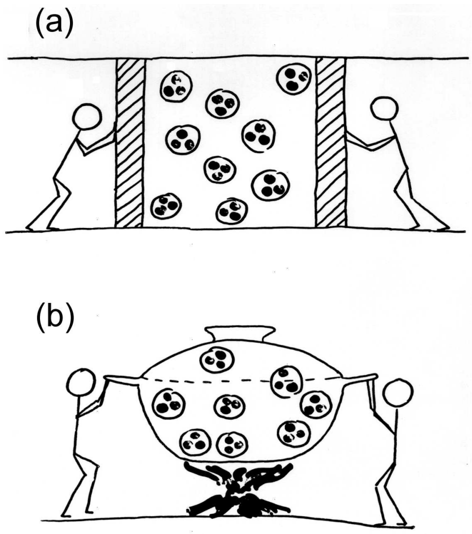

Fig. 3. Two technologies to produce the dense hadronic matter: compression (a) and heating (b).

Let us now consider the gas of noninteracting pions. Its temperature is assumed to be so high that the pions can be treated as massless particles. (The approximation appears to work quite well even for the temperatures which are close to the pion mass equal 140 MeV.) The pion density is then

$$\rho_\pi = \int \frac{d^3p}{(2\pi)^3} \frac{g_\pi}{\mathrm{e}^{E/T} - 1} = \frac{g_\pi \zeta(3)}{\pi^2}\, T^3 \cong 0.73\, T^3 \,, \tag{3}$$

where $E \equiv |\mathbf{p}|$; $\zeta(z)$ is the Riemann function and $\zeta(3) \cong 1.202$; g_π is the number of the particle internal degrees of freedom, which is 3 for the gas of π^+, π^0 and π^-. One immediately finds from (3) that the inter-pion separation is smaller than 1 fm for $T > 219$ MeV.

The deconfinement phase transition has been studied in the numerous phenomenological approaches. We consider here the simplest possible model, where the phase transition is assumed to be of the first order. Then, one can apply the Gibbs criterion to construct the phase diagram. The hadron phase is modeled by the ideal gas of massless pions of three species ($g_\pi = 3$) while the quark–gluon one by the ideal gas of quarks and gluons, which is baryonless *i.e.* the total baryon charge vanishes due to the equal number of quarks and antiquarks. Let us compute the number of the internal degrees of freedom in the quark–gluon gas. One should distinguish here the fermionic degrees of freedom of quarks g_q and the bosonic of gluons g_g. There are two light quark flavours (u and d); there are quarks and antiquarks; we have two spin and three color quark states. So, $g_q = 2 \times 2 \times 2 \times 3 = 24$. The gluons are in two spin and eight color states, which give $g_g = 2 \times 8 = 16$.

Since the pressure of the ideal gas of massless particles equals one third of the energy density, the pressure exerted by the pion gas is

$$p_\pi = \frac{1}{3} \int \frac{d^3p\, E}{(2\pi)^3} \frac{g_\pi}{\mathrm{e}^{E/T} - 1} = \frac{g_\pi \pi^2}{90}\, T^4 \cong 0.33\, T^4 \,, \tag{4}$$

while that of the quark–gluon plasma equals

$$p_{qg} = \frac{1}{3} \int \frac{d^3p\, E}{(2\pi)^3} \left[\frac{g_g}{\mathrm{e}^{E/T} - 1} + \frac{g_q}{\mathrm{e}^{E/T} + 1} \right] = \left(g_g + \frac{7}{8} g_q \right) \frac{\pi^2}{90} T^4 \cong 4.1 T^4 \,. \tag{5}$$

According to the Gibbs criterion the phase, which generates the higher pressure at a given temperature, is realized. Then, one finds from eqs. (4, 5) that the pressure of the quark–gluon plasma is always greater than that of pions. Therefore, we should have, in conflict with the experiment, the quark–gluon phase at any temperature. However, we have not taken into account the pressure exerted by the vacuum on quarks and gluons. Subtracting the bag constant B from the r.h.s. of eq. (5), one finds that below the critical temperature T_c there is the pion gas and above the quark–gluon plasma. The critical temperature is given as

$$T_c = \left[\frac{90B}{\pi^2 (g_g + \frac{7}{8} g_q - g_\pi)} \right]^{1/4} \cong 0.72\, B^{1/4} \,.$$

Taking $B^{1/4} = 200$ MeV, we get $T_c = 144$ MeV.

In Fig. 4 we show a schematic phase diagram of the strongly interacting matter. The baryon density is measured in the units of the normal nuclear one. The point at $\rho = \rho_0$ and $T = 0$ represents normal nuclei. In fact, this is the only point in the diagram which is really well known and understood. At $\rho > \rho_0$ and $T = 0$ there is a region of the dense nuclear matter with several exotic forms being suggested. When the baryon density exceeds $2 - 3\rho_0$ one expects a transition to the quark–gluon phase. At even higher densities the plasma is believed to be perturbative *i.e.* weakly interacting. So, one deals with the quasi ideal strongly degenerated quark gas. The nuclear matter at the temperature larger than a few MeV is traditionally called a hadron gas being mostly composed of pions when the baryon density vanishes. Then, we have quite reliable QCD lattice results which tell us that there is the deconfinement phase transition at $T \cong 180$ MeV. With the significantly larger temperatures we approach again the perturbative regime where the plasma is weakly interacting. In the next section we discuss how the phase diagram can be explored by means of the heavy-ion collisions.

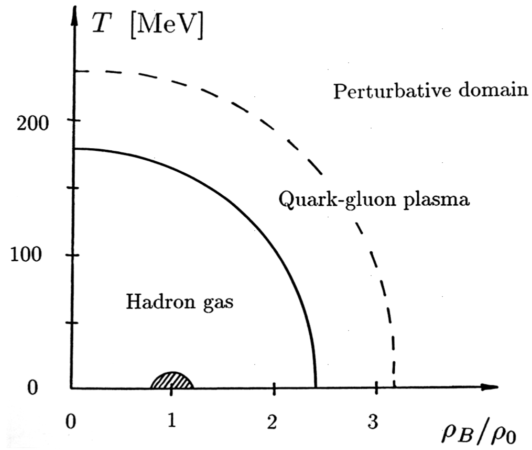

Fig. 4. The phase diagram of the strongly interacting matter

6. Heavy-ion collisions

As already mentioned in the introduction, the heavy-ion collisions provide a unique opportunity to study the quark–gluon plasma in the laboratory experiments. More precisely, a drop of a dense hadronic matter can be created in the collisions; the question whether the matter is for some time in the deconfined phase is a separate issue.

The physics of heavy-ion collisions essentially depends on the collision energy. A good measure is not the whole energy of the incoming nucleus but the energy per nucleon. The point is that at the energy of a few GeV per nucleon, which is the lowest energy being interesting from the point of view of the quark–gluon plasma, the nucleus does not interact as a whole but there is an interaction of the overlapping parts of the colliding nuclei as depicted in Fig. 5. The nucleons to be found in these parts are called the *participants* while the remaining ones the *spectators*. One often distinguishes the central from the peripheral collisions. In the latter ones, which proceed with a large value of the impact parameter, most of the nucleons are spectators. In the case of the central collisions basically all nucleons from the smaller nucleus (target or projectile) are participants and the interaction zone is the largest. Obviously the central collisions are the most interesting in the context of the quark–gluon plasma searches. Unfortunately, the nucleus–nucleus cross section is dominated by the peripheral collisions. The cross section contribution of the collisions with a given impact parameter b is $2\pi b db$. Therefore, the contribution vanishes when $b \to 0$.

The high-energy nucleus–nucleus collision proceeds according to the following scenario. The overlapping parts of the colliding nuclei strongly in-

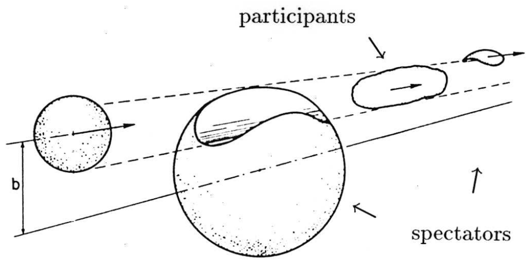

Fig. 5. The geometry of the nucleus–nucleus collision at high energy

teract and a dense and hot hadronic system, which is often called a *fireball*, is created. Initially it is formed either by the quarks and gluons or by the hadrons. Then, the system expands and cools down. If the fireball matter has been initially in the deconfined phase, the system experiences the hadronization *i.e.* the quarks and gluons are converted into the hadrons. The fireball further expands. When its density is so low that the mean free path of the hadrons is close to the system size or the expansion velocity is comparable the that of the individual particles, the whole system decouples into the hadrons which do not interact with each other any longer. However, the unstable particles can still gradually decay into the final state hadrons. The moment of decoupling is called the *freeze-out*, because the particle momenta are basically fixed at that time. Thus, the final state hadrons, which are observed by experimentalists in the particle detectors, characterize the fireball at the freeze-out.

As we discussed in Sec. 5, the dense hadronic matter can be obtained either due to the nuclear matter compression or the nuclear matter heating. It appears that the significant effect of compression at the early stage of heavy-ion collisions, which is still insufficient for the quark–gluon generation, is observed at the relatively low energies, not larger than 1 GeV. Then, one can study the dense nuclear matter of rather small temperature. At higher energies the atomic nuclei appear to be rather transparent *i.e.* the colliding nuclei traverse each other. The participants are only slightly deflected from their straight line trajectories due to the interaction. However, they lose a sizable portion of their energy which is further manifested in the form of produced particles, mainly pions. Therefore, the baryon density of the produced hadronic system is not much increased even at the early collisions stage. When the produced system expands, the baryon density soon gets a value which is smaller than the normal nuclear density. The strongly interacting matter however is significantly heated up. If the temperature exceeds the deconfinement transition temperature, the matter is expected to be in the quark–gluon phase.

A comment is in order here. We use the concept of the temperature which explicitly assumes that the system is in the thermodynamic equilibrium. It is far not obvious that this is really the case. Any physical system needs some time to reach the state of equilibrium. The hadronic matter produced in heavy-ion collisions cannot achieve the *global* equilibrium but the theoretical as well as experimental arguments suggest that the *local* quasi equilibrium is possible. The global equilibrium is characterized by the thermodynamic parameters which are unique for the whole system. Before the global equilibrium is achieved a system is usually for some time in the local equilibrium state. The system's parts are then already in the equilibrium but the thermodynamic parameters — temperature, density, hydrodynamic velocity — vary from part to part. The hadronic matter produced in heavy-ion collisions is however not kept in any container but it immediately starts to expand, mainly along the beam axis. Consequently there is a sizable variation of the hydrodynamic velocity in this direction.

7. Plasma signatures

The plasma creation is expected to occur at the early collision stage, but it must hadronize *i.e.* experience the transition to the hadron gas, when the matter is expanding and cooling down at the later stages. Therefore, we always observe hadrons in the final state of the collisions and it is really difficult to judge whether the plasma has been present or not. Although it is hard to imagine a smoking gun proof, a few signatures of the plasma creation has been proposed. We discuss below the two which seem to be the most promising.

It has been argued that the presence of the plasma at the early collision stage increases the number of strange particles observed in the final state. The point is that the mass of strange quark mass appears to be significantly smaller than that of strange particles. The strange (s) quark mass is about 150 MeV. Therefore, one needs 300 MeV to produce the $s - \bar{s}$ pair. The strange quarks must be, of course, produced in pairs because the strangeness is conserved in the strong interactions. The most energetically favorable way to produce the strangeness at the hadron level proceeds in the reaction $\pi + N \rightarrow K + \mathit{\Lambda}$. Here the incoming particles — the pion and nucleon — must curry over 500 MeV in the center of mass frame. Thus, it is easier to produce strangeness at the quark–gluon than hadron level. Once the strange quarks appear in the plasma, they cannot disappear — a rather rare annihilation process of $s - \bar{s}$ pairs can be neglected — and consequently, they are distributed among the final state hadrons. Quantitative comparison of the two scenarios with and without the plasma indeed shows that the presence of the plasma leads to the significant strangeness enhancement.

The second plasma signatures deals with the J/ψ particle or charmonium which is a bound state of the charm quark c and antiquark $\bar{c}$. The J/ψ particle is expected to dissolve much easier in the quark–gluon environment than in the hadron one. This can be understood as a result of screening of the potential, which binds the c and $\bar{c}$ quarks, by the color charges of the plasma particles. Therefore, the number of the J/ψ particles in the final state should be significantly reduced if the plasma is present at the collision early stage.

The two predicted quark–gluon plasma signatures have been indeed experimentally observed in the central heavy-ion collisions at the energy of about 200 GeV per nucleon which have been recently studied at CERN. The whole set of the experimental data however does not fit to the theoretical expectations. There have been advocated the mechanisms of the strangeness enhancement and J/ψ suppression which are different than those mentioned above. Therefore, it is a matter of hot debate whether the plasma is produced at the energies which are currently available. The plasma generation at higher energies seems to be guaranteed.

8. Perspectives

In the near future the nucleus–nucleus collisions will be studied at the accelerators of a new generation: Relativistic Heavy-Ion Collider (RHIC) at Brookhaven and Large Hadron Collider (LHC) at CERN. The collision energy will be larger by one or two orders of magnitude than that of the currently operating machines. The heavy-ion experiments have been performed till now in the beam-targets system where the accelerated ion is smashed against the target nucleus which is initially at rest. RHIC and LHC will use another principle — there will be accelerated two intersecting ion beams. The energy of each beam will be 100 GeV per nucleon at RHIC and 3 000 GeV at LHC. Thus, the collision energy in the center of mass frame will equal, respectively, 200 and 6 000 GeV which should be compared to 20 GeV presently available in the beam-target systems.

The proton–proton collisions have been already studied experimentally at the energy domain of RHIC. The perturbative QCD, which is hardly available at lower energies, is extensively used here. This is possible because the average momentum transfer grows with the collision energy and the QCD coupling constant is then relatively small. Therefore, the theoretical understanding of the collisions, paradoxically, improves with the growing energy.

The nuclear collisions at RHIC or LHC are expected to be so violent that the quarks and gluons comprising the nucleons will be easily deconfined, for some time of course. At such huge energies the nucleon is visualized as a cloud of partons which breaks up into the parton showers as a result of collision with another nucleon. Thus, the creation of the quark–gluon plasma seems to be unavoidable at RHIC and LHC. The method of its detection however remains to a large extend an open question.

I am very grateful to Marek Gaździcki for critical reading of the manuscript.

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